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Editors

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Controlled Nucleosynthesis

Breakthroughs in Experiment and Theory



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Controlled Nucleosynthesis

Fundamental Theories of Physics

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Controlled Nucleosynthesis

Breakthroughs in Experiment and Theory

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INTRODUCTION

This collection of papers presents the main results of a logical analysis of extensive theoretical and experimental searches for efficient ways to overcome two long-standing scientific and technical problems: (1) the creation of a driver of inertial thermonuclear synthesis with a positive gain of energy and (2) the neutralization of artificially created radioactive materials accumulated as a product of the human activity.

The collection reports the main achievements of theoretical and experimental studies performed by a group of experts in a number of relevant areas: the synthesis of stable dynamical structures of different physical natures; focusing and self-focusing of electron and plasma beams; concentration and self-concentration of energy in material media and physical structures on different scales; and the study of both the nuclear combustion mechanism of substances under laboratory conditions and the chemical elements produced by such a combustion.

The investigations were carried out in a framework provided by a privately financed program “Luch” designed to solve the problem of finding an efficient and safe technology for pulse initiation of controlled nuclear combustion at the Electrodynamics Laboratory “Proton-21” (Kiev, Ukraine) during 2000–2004.

The papers of this collection are written by the immediate authors of new ideas and constructions, designers of methods, and executors of physical experiments and measurements. These professionals had the good luck to solve the problem of initiating the controlled nuclear combustion of substances without an accompanying creation of radioactive “ashes.” By taking their clues from the processes occurring in exploding stars, they discovered and, in first approximation, investigated a great number of physical processes and phenomena which were not predicted in most cases, being deemed improbable, or were born only in the minds of romantic visionaries.

Most of the artificially initiated and first-studied processes and phenomena (the full-scale laboratory nucleosynthesis of stable light, medium, heavy, and superheavy nuclei and atoms; the neutralization of radioactivity in the process of pycnonuclear combustion under conditions of the self-development of an artificially initiated collapse of substance; the low-energy induction of the decay of the unknown stable superheavy nuclear structures [which are, possibly, nuclei] surrounded by “classical” earth-related substances), including those which have been reproduced hundreds or thousands of times in our Laboratory for five recent years, remain unperceived

and frequently rejected, in principle, by many members of the physical community. Such a state of affairs has many objective reasons and explanations.

Among them are: the fact that both our derived results and their variety seemed inexplicable at first glance; the apparent conflict of the announced results with tradiational ideas on the possible conditions and mechanisms governing nuclear reactions; the existing absence of official reproduction of our main results by other groups and laboratories, which shortcoming is due to the incompleteness of patenting procedures and the commercial character of the project, which factors often bedevil this sphere of activity; and the absence of a commonly known and accepted physical theory, from the perspective of which one can explain the totality of data reported by the Proton-21 group.

However, the main reason for the distrust and rejection of our results, which is really unexpected in many respects and sensational in a certain sense (first of all, from the viewpoint of the potentialities they open up for an *explosive stage of development in our technological civilization in harmony with terrestrial and cosmic realities*), is a deficit of detailed and systematically presented information on: the underlying ideas; schemes of performed physical experiments; procedures of investigations; measuring tools; real totality of experimental data, which look fantastic but are quite real; reproducibility of observed effects; and the great number of details and nuances composing the essence of any physical experiment and allowing the objective estimation of its certainty.

The aim of the present work is to partially fill in the aforementioned gaps and to give general information on:

- the variational principles and conceptual models forming the basis of the original (and, possibly, unique) means of the artificial initiation of the collapse of energy-concentrating targets;
- the organization of the first series of successful experiments concerning the artificial initiation of a self-supporting pulse process of nuclear combustion;
- the results of measurements and estimates of the parameters of an electromagnetic driver that induces the coherent quasi-isentropic shock compression of substances;
- the preliminary results of studying the dependence of the abundance of the observed chemical elements, constituting the products of the nuclear regeneration of the initial substance of targets, on the main nuclear-physical characteristics of the process;
- the peculiarities of the emission spectra of the “hot dot” of a collapse in the entire ranges of wavelengths of the electromagnetic emission and

of masses and energies of the particles forming the shock front of an explosive plasma bunch ejected by an exploding target;

- the theoretical foundations and the physical models of reproducibly observed unconventional phenomena.

We thank, in anticipation, all the readers who will make remarks on the substance of the presented results, which are no doubt far from perfect, and especially those researchers who will verify our results with the use of procedures overlooked by us or with technical facilities inaccessible to us. We have little doubt that readers who acquaint themselves with the process discovered by us, which is fantastic and inexhaustible in its manifestations, will be overwhelmed, as we have been, by its harmony, power, and perfection.

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My coauthors and I express our gratitude and respect to the investors in our project, “Proton-21” shareholders, and “PrivatBank” owners Gennadiy Bogolubov and Igor Kolomoysky, who manifested an outstanding mental outlook and foresight. Without the financial support of these individuals, the outcome of our scientific project would obviously have been gravely imperiled. My friendly gratitude goes also to Igor Didenko, who entered our project in the final stage of book preparation.

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Stanislav V. Adamenko

Part I

Approach to the Problem

1

PREHISTORY

S. V. Adamenko

With deep gratitude to my father.

At the beginning of 2003, Professor Yurii Kondrat'ev got to know the results derived at the Electrodynamics Laboratory "Proton-21" and then gave an account of his impressions to Professor Franco Selleri. In autumn of that year, when Selleri visited the NASU Institute of Mathematics in Kiev at the invitation of Kondrat'ev to give a lecture, he also was our guest for several days. I had the pleasure to show him the laboratory's facilities and to tell about our experiments, our ideas about the mechanisms underlying the astonishing physical phenomena discovered by us, and the bases of our assertions about their existence in nature, in general, and their reproduction in our laboratory, in particular.

Selleri readily comprehended the difficulties we had encountered when trying to publish the results of our experiments on the initiation of nuclear combustion and laboratory nucleosynthesis in refereed journals. In the great majority of cases, the conclusions of referees consisted literally of several phrases which were based on three fundamental, in their opinions, positions:

1. This cannot occur in principle; the assertions of the authors about the controlled realization of collective nuclear reactions in a superdense substance are based, most probably, on the incorrect interpretation of the results of measurements.
2. The experimental results declared by the authors have no theoretical substantiation and contradict established physical ideas.
3. The authors propose the theoretical models of nonexistent physical processes.

The recommendation of Selleri was a very constructive one: "In cases similar to yours, it is very difficult to destroy the wall of distrust by piecemeal publication of the papers devoted to separate aspects of the project. I think it is necessary to quickly prepare an anthology of papers which must include the most important things, starting from the conception of the experiments

and finishing with the presentation of the proposed theoretical models and mechanisms of the discovered phenomena.”

Appealing to me personally, Selleri added: “You should also not forget to tell the history that led you to these problems, i.e., when and why did you become interested in nuclear synthesis?”

In this context, the dedication to my father prefacing this chapter is not the usual expression of filial appreciation. Indeed, if my father held a pedestrian view of life and parental obligations, I would have no special reason for evoking his memory in order to explain why I became motivated to tackle a purely physical problem from the traditional viewpoint, not being a professional physicist myself.

My father was an extraordinary person in many ways. In particular, he had a phenomenal memory that enabled him to recall and use, at any moment and over many years if necessary, an inconceivable, from my viewpoint, number of dates, names, poems, quotations, facts of the own life, etc. This excellent memory and the ability to read rapidly caused my father to become an erudite person. He was especially interested in scientific and technical novelties and achievements, reports about which were numerous in the 1950s and 1960s.

From childhood, he dreamed about becoming a medical doctor. But in 1939, at the age of 17, he was called up for the military service in the Soviet Army. Then, for the first 20 years of his long-term service, he tried many times, but without success, to go into retirement or, at least, to get permission to enter the military-medical academy, which was far removed from his military profession. Recalling the imaginative mind-set my father revealed in the process of my upbringing and the adult role games he invented for me and my friends, I am sure that he was also a real teacher at heart.

When I was in my fourth year, my father apparently thought it was time to teach me the virtues of work and having a purpose in life. He brought home a large ball bearing and challenged me to extract smaller balls from it. I remember well how acutely I wanted to get them by myself and how simple the problem seemed at first. However, the ball bearing resisted my initial efforts, and smaller balls did not jump out themselves. I had to grab a file and begin to work. I do not recall how long this went on, but only remember that this Sisyphean labor annoyed me only when I realized that I would be filing a long time, at least several days. So I complained to my father. He was quick to advise that difficult tasks should be solved first in one’s head. Only if the solution becomes clear, a hand may reach for a tool. As an example, he told me about tricks a monkey had to use in order to access food frustrating conditions.

A prompt helped me. I understood that, firstly, the ball bearing will be split up if it is thrown onto a stony roadway. Secondly, the resulting fragments will not disperse if the ball bearing is first placed in a small knotted bag. We together executed the experiment, and the fragments were in my hands in no time.

This was my first creative success preserved in my memory as an example of the efficiency resulting from a proper approach.

Seven years later, during an evening walk with my father, I was given a task whose comparatively simple solution I searched for most of my life. It was November 1, 1958; I remember the date only because I was ten the next day. The main theme of our conversation was that, at that age, it was time to think about serious matters and to prepare for adult life, rather than to squander free time without any purpose in mind.

We looked at the evening sky, and my father taught me how to find the Polar star and the easily recognized constellations. See, he said, stars differ in brightness and even in color, because they are at various distances from us and have different sizes and temperatures. But they are all similar in principle to our Sun. Stars are shining very long, for billions of years. Then they become dim and collapse. Further, some stars explode. The radiant energy of stars originates in the combustion of matter. But it is not ordinary chemical combustion, like that in a campfire, but rather a thermonuclear one, wherein the lightest nuclei of hydrogen form the nuclei of heavier chemical elements by fusing many times. Physicists name processes of this kind thermonuclear synthesis.

In thermonuclear fusion, the amount of the released energy is millions of times that produced in the usual combustion of coal or gas. The fusion process was already realized on earth in the explosions of hydrogen bombs. If the energy of such explosions were to be used for peaceful purposes, the demands of humans for energy would be satisfied for thousands of years.

Unfortunately, this prospect is presently out of reach—for the following reason: A thermonuclear charge can now be fired only by the explosion of an atomic bomb, for which a critical mass is required. Thus, an atomic bomb cannot be made so small that it does not destroy everything for tens of kilometers around it. Consequently, scientists are now faced with the problem of inventing a trigger for thermonuclear charges that is simpler and cheaper, so that it can be permanently used in a thermonuclear reactor producing heat and electricity.

It turns out that this problem is incredibly difficult and expensive to solve. Scientists from various countries have tried to solve it jointly. If one is interested in it, one can become a physicist and possibly, devise a suitable solution.

- But why has this problem not been solved already, and what must be done?
- Well, it is necessary to heat hydrogen to an extremely high temperature, much higher than that of the sun; and no such technology exists at present.
- But we can place hydrogen at the focus of a great magnifying lens and heat it in such a way to any temperature!
- This method leads nowhere.
- Why?
- Things are not as simple, as it seems. The mastering of such a source of energy is a very complicated problem, though the experts believe that the problem is not hopeless. Learn, examine, and dream! Anybody has a chance if he or she tries. As is known, complex problems sometimes have simple solutions.

I often recall the evening conversation with my father about stars and the tempting subject of nuclear synthesis as a particularly seminal event of my childhood.

In the years that followed, the problem he first posed attracted me more and more. I can give no rational explanation why the persistent thoughts about the possible, from my viewpoint, mechanisms and nature of nuclear synthesis became a habit, a hobby, as it were—one that did not require separate time, since it settled in the back of my mind, where it nonetheless kept my imagination in training.

For many years, I had no serious plans for solving the synthesis problem, as I could not imagine that my own contribution would be very meaningful in comparison with the efforts of true experts. So my musings remained on an amateurish level.

Considering stably functioning biological and technological systems composed of simple elements of the same type, I searched for a hidden logic of their origin and evolution, and for a mechanism that could help me with the unsolved problem of controlled nuclear synthesis.

I was troubled by the fact that a process that produces enormous amounts of energy, serves as its main source in the universe, arises in stars spontaneously and keeps running there by itself, does not seem reproducible under terrestrial conditions, despite all our scientific and financial efforts.

For a long time, it seems illogical that the simplest energy process of fusion of light nuclei, in which cosmic nuclei easily participate for objective reasons, is not realized in a controlled and efficient manner on earth, even though scientists have gained a high degree of comprehension about the structure and behavior of atomic nuclei as well as the formalization and

mathematical description of the processes involving them. I got the impression that the search was carried out in the wrong direction and that efforts were aimed at forcing nuclei meant to be fused into a behavior not peculiar to them, though fusion could occur with the selection of a more favorable final state.

This simple thought served as a first prompt: It is necessary to abstract from a seemingly inevitable *final state* of a system of nuclei and to consider more attentively the possibly optimum *conditions for the transition* of this system from a given initial state to the required final one. While pursuing postgraduate studies and learning the mathematical methods for the synthesis of multiply connected dynamical systems with optimum stability, I undertook a search for, and an analysis of, such regularities in the synthesis of the complex systems composed of interacting (exchanging energy or information) elements of any physical nature which would be common for nuclear and, e.g., biological or controlling structures. The main question was as follows: *For what* reason do “independent” elements combine to form systems and ensembles restricting their freedom? What conditions control the sizes of systems, the number of primary elements, and the structure and the force (energy) of bonds between? What criterion guides the initial building blocks into forming a particular final product? Is it possible that, in order to complete this quest, they have randomly to “survey” all possible structural variations, among which would appear the unique required solution? Biology and genetics answer the last question in the negative: Admissible solutions would have caused *Homo sapiens*, who invents questions and searches for answers, *not* to appear for still many billions of years! Changing structure during their development, complex self-organizing systems progress to their optimum structure along a route that differs slightly from the shortest possible one. This can only mean that the systems are surely led by some force. In this case, in each step of their development (rearrangement), a systems can “estimate”, in some manner, the degree of its imperfection and can “detect” the reason for it, by getting a stimulus for the next step on a gradual track to the optimum state. It is well known that every stable system is certainly optimum in some sense. In theory, the corresponding criterion of optimality can always be found on the whole by solving the inverse problem of synthesis for the system under study, such as a traffic network, living cell, atomic nucleus, or atom. Every self-organizing structure has its own system-forming criterion. By what does it differ from a set of other possible ones, what does it demand from the system, how does it appear, and how does it take into account the features of the system and the conditions defining its self-development? In the mid-1970s, while engaged in my postgraduate tasks, I searched for an analytic solution to the problem dealing with the synthesis

of an optimum control with feedbacks for a controlled linear dynamical system subject to restrictions in the form of proper (invariant) linear subspaces that are specified *a priori*, in other words, with restrictions on the phase portrait of the optimum multidimensional dynamical system. As a tool, I at first used the classical method of analytic construction of optimum regulators (see Refs. 1–3). But it rapidly became obvious that no analytic methods for the systems with the mentioned restrictions existed. The development of appropriate methods was the theme of my dissertation. In particular, I proposed the *method of binary synthesis*, whose peculiarity resided in the following: Contrary to the classical approach, the quality criterion for a transient process as a measure of integral excitation of a system in the spaces of phase coordinates and controlling actions was set in the form of an integral of the so-called optimum target function, rather than by the integral of the sum of *a priori* positive-definite functions of the phase coordinates and controlling actions. This last sum becomes the optimum target function (but *a posteriori*!) when the unknown control vector $\mathbf{U}(t)$ as a function of time is replaced in the corresponding term of the sum by the required optimum law of control with a feedback $\mathbf{U}^0[x(t)]$, i.e., we have a function of phase coordinates of the optimized system.

The sense of such a transformation of the classical problem of the synthesis of an optimum dynamical system (it turns out to be a generalization) consists in the fact that the *a posteriori* optimum target function is the Lyapunov function for an optimized closed system (a positive-definite quadratic form of the phase coordinates in the linear case), whose set of quickest-descent trajectories approaches (as closely as desired with a weakening of the restrictions on the control) the phase portrait of the optimized system. Thus, the *a priori* setting of an optimum target function allows one to form any desired phase portrait in the process of synthesizing the optimum system, i.e., the eigenfunctions of the dynamical system and its eigennumbers.

Application of the method of binary synthesis to optimize dynamical systems allowed me to get a number of interesting results.

The first formal result consists in the conclusion that one can completely reject the necessity, inevitable in the classical case, to seek the parameters of the *a priori* quadratic forms of a quality functional (the criterion of optimality) by the method of an actually arbitrary exhaustive search in order to get the more or less satisfactory phase portrait of an optimum system. I note that such a portrait can nevertheless never “fall” into the set restrictions. Instead, the *necessary*, but basically *not guessed*, parameters of the *a priori* quadratic form of the criterion and parameters of the optimum feedback law were finally derived in the framework of the method of binary

synthesis from the solutions of the relevant systems of equations—which are, naturally, also the generalizations of the systems of equations belonging to the classical version.

In other words, two optimization problems were solved within the framework of the method of binary synthesis, at least from the formal viewpoint. Our direct problem was to search for $\mathbf{U}^0[x(t)]$, while the inverse problem consisted in the search for a term in the integrand of the quality criterion as an *a priori* indefinable quadratic form in the phase coordinates of the system. This term is *a priori* unknown but uniquely necessary for the application of the required restrictions. Just this circumstance explains our use the adjective “binary” to qualify the proposed method of synthesis.

The second, more significant, result of this method is that its natural motions relative to, e.g., any hyperplane restriction given in phase space can possess, if necessary, an arbitrarily small inertia. Moreover, the given hyperplane-restriction can differ slightly, to any desired degree, from the $(n - 1)$ -dimensional (n being the dimensionality of phase space) proper invariant subspace of the optimized system. In particular, this means the minimization of the excited (forced) motion energy of the optimized system relative to the own relevant hyperplane arbitrarily located, in the general case, in phase space.

It is easy to see that this actually implies that the procedure of binary synthesis allows one to attain the maximum stability and minimum dissipativity of the system with respect to a “pathological” external perturbation which moves the point representing the system beyond the given invariant hyperplane possessing the highest priority among the goals of the optimization or those of the homeostasis of the system.

The external perturbation setting a direction in the system’s phase space such that the forced movement along it is characterized by the maximum absorption of the energy dissipated in the system, which leads to its maximum heating or to the maximum destruction, was called the “dominant perturbation”. A dominant perturbation exciting synchronously and with identical phase all the degrees of freedom or all interacting elements of the system is called the “global dominant perturbation”.

It was easy to see that, in the framework of the problem of binary synthesis, the optimum structure of a multiply connected dynamical system and parameters of the formal criterion can be found in a self-consistent way as functions of the exceptionally objective factors: namely, the parameters of an object of the optimization and the parameters of external perturbations acting on it.

It thus turned out that, in the problem of binary synthesis under the condition of the setting (or the determination by the system itself) of the

direction of action of the global dominant perturbation, the determination of an optimum structure of the bonds between elements of the system can occur self-consistently and simultaneously with the search for the corresponding (not fixed *a priori*) parameters of the criterion of optimality of the process. Similar to what holds for the optimum system, this process depends on its initial parameters and also on external perturbations acting on the system. Upon a change in external perturbations, we can observe, in principle, the automatic “switching on” of the next cycles of the adaptation of the system to external conditions. These cycles, by repeating with each recurrence of a dominant perturbation, are able to support the process of continuous self-organization and reorganization of the system and to ensure asymptotically, in particular:

- The “reflection of images” of the external dominant perturbations in the structure of bonds between elements of the system and, as a consequence, in the set of its own subspaces of different dimensionalities, which can be considered as the operation of a distinctive mechanism of the system’s memory and its adaptation to the external dominant perturbations.
- The maximization of stability and the minimization of inertia for the natural motions of the system that arise in its phase space as a result of the action of external dominant perturbations.

The statement of the mentioned peculiarities of the problem and the binary synthesis algorithm for a dynamical system made a strong impression on me in 1980. In these peculiarities, one may guess the characteristic features of a long-expected mechanism of the “self-synthesis” to lie. The last was comprehended as the self-organizing self-developing “not powerful” natural process of nucleosynthesis, temporarily unclarified, but certainly existing and held responsible for the formation of a whole set of the naturally coexistent nuclei and atoms of chemical elements.

Against the background of the discovered peculiarities of the offered algorithm of the optimization of structures, I made an assumption on the existence of a universal natural regularity which I called tentatively the *principle of regularization of perturbations and dynamical harmonization of systems*. This regularity indicates the general direction for the improvement of self-organizing multicomponent dynamical systems: At the expense of a restriction of the individual freedom of interacting elements (particles), one can reach a maximally attainable decrease in the inertia of a reaction of the whole system to various external dominant perturbation that coherently act on each participating element and thus have the distinctive signs of a mass force.

At that time, such a “discovery for internal use” caused me to experience an intense emotional excitement. I remember well-being overcome by a spell of euphoria.

Thus, the process of synthesis of self-organizing dynamical systems, which one can realistically apply to nuclear structures as well, reveals the logic of initiation and development that appeals to my way of thinking.

Intuition prompted me to surmise that, regardless of the degree of practical usefulness and real novelty of the formulated principle, the apparently self-sufficient physical mechanism of reflection on the information contained in the structures of dynamical systems could become a peculiar key to comprehending the required self-organizing mechanism of nuclear synthesis on the macroscopic level, which continued to be a castle in the air. Analysis of a system optimized in the framework of the binary synthesis algorithm showed that a parallel consequence of such a “behavior” of the self-organizing dynamical system will be a maximization of the stability of its own motion excited by a reflected external dominant perturbation, as well as the minimization of a “destructibility” of the system under the action of this perturbation; this can be interpreted naturally as the maximization of the binding energy of the system with regard to restrictions on the physical nature of system-forming elements and the forces of their mutual interaction.

At that time, the computer realization of the binary synthesis algorithm showed that it is possible to attain an arbitrarily small inertia of the system in the direction of the action of a dominant external perturbation upon a sufficiently large number of “bound” (interacting) elements, despite a restriction on the forces of interaction (on the intensity of bonds) between them; in this case, the system’s inertia on each of the remaining degrees of freedom can be arbitrarily slightly different from the initial one.

The mentioned peculiarities of the organization of optimum systems, despite the colossal differences of the used limitedly simplified descriptions from adequate physical models of atomic-nuclear and other natural structures, allowed me to assume that the basic synergetic properties of systems did not depend on their specific nature and always manifest themselves in the self-organization of complex dynamical structures undergoing the organizing action of intense (dominant) external perturbations.

By the beginning of 1986, similar thoughts led me to conclude that the main difficulty for the controlled synthesis of nuclei consisted in the artificial creation of just such external dominant perturbation common to the totality of nuclei involved in the process of synthesis, whose optimum “reflection” would be completed, in the sense of the above-presented approach (the binary synthesis), by the exothermic (exoenergetic) fusion of initial nuclei.

I further reasoned as follows:

If the factor defining the result of any synthesis is a dominant perturbation common to the interacting initial components, such a perturbation must exist and play a defining role in the natural processes such as a simple fusion of nuclei (i.e., in the “thermonuclear synthesis” in the traditional sense) and, in the wider sense, the complex process of natural nucleosynthesis, whose products are the nuclei of all chemical elements up to the heaviest ones. If such system-forming action were to be discovered, then it would be possible to search for its analog artificially realized under laboratory conditions.

These were the preliminary positions, in general terms, of my conception regarding the artificial initiation of self-organizing nuclear synthesis, which were formulated in 1987–1988, 30 years after I first encountered the problem.

If it were not for an improbable coincidence of circumstances, the necessity or even the occasion to tell this history would never have occurred. At least, I had no such intentions until February 2000.

At that time, a decade after the start of the “Perestroika” in the USSR and five years after Ukraine gained its independence, I together with many colleagues in the profession had to work in the field of business and already saw the decline of personal dreams about thermonuclear synthesis. Moreover, the inexorable chain of some events deprived me of the last hope to be involved the solution of the thermonuclear synthesis problem. But the situation was about to change due to unforeseen circumstances.

In 1996, my dear friend Dr. Boris Sinyuta, an expert in the field of radiation medicine, who since passed away to my deep regret, introduced me to Dr. Vladimir Stratienko of the Kharkiv Physico-Technical Institute in order to discuss commercial plans for the production of isotopes for medical purposes in Ukraine.

Dr. Stratienko saw me as a former young scientist and now a businessman who, on the one hand, had some money and, on the other hand, was ready to share it for the benefit of nuclear science and scientists, if an interesting project presented itself.

So, Dr. Stratienko tried to convince me that, by using microbeams of relativistic electrons in a vacuum diode, it was possible to focus them at the end of a thin cylindrical target anode up to a current density of at least $10^{10} \text{ A} \cdot \text{cm}^{-2}$. He assumed that this would lead to the formation of a highly ionized plasma in a small near-axis volume where the beam interact with the target. Such a plasma could be compressed and held by the magnetic field of the beam in a state with density and temperature high enough to ensure

a positive energy-gain in the scheme of inertial thermonuclear synthesis if a suitable thermonuclear (e.g., D-T) target is used.

To practically design a driver on the basis of a hypothetically existent mechanism for the self-focusing of an electron beam to extreme current densities, a relatively small financial support by a group of enthusiasts, living in the difficult period of the global economic transformation, was required.

At that time, there was not, and could not be, any convincing proof of the feasibility of the proposed scheme. However, all doubts were rejected on the basis of a presentiment, rather than of a comprehension, that a focused electron beam can fully play the role of a long-expected dominant perturbation for the set of particles forming the completely ionized substance of a target and compressed by the magnetic field of the beam. Strangely, an inner voice commanded us to act, promising the realization of a dream in the face of an adverse reality.

The temptation was great, and the sum of money required for support of the pilot was small and available. So, the decision was made, business was ceased, and I began recounting time in my last attempt to participate in solving the problem of controlled thermonuclear synthesis. Soon an initiative group, composed of Kiev and Kharkiv experts (mainly theorists) in the fields of solid state physics, plasma physics, high-energy physics, accelerating techniques, the theory of systems, and nuclear physics, was formed.

For at least three years, we held on a frequent basis seminars in Kiev and Kharkiv in turn; we discussed mechanisms, models, analogies, theories, the experiments performed by others, and plans for the establishment of a laboratory, in which we hoped to solve, by simple and efficient means, the problem of controlled thermonuclear synthesis in its inertial version. In the first stage, we presume to achieve success with the help of a self-compressing, “self-lacing” hard-current microbeam of electrons. The beam directed on the end or point of a target-anode should move continuously along the target axis, “consume” its core, and transform it into a superheated, supercompressed thermonuclear plasma until the pulse ceases. This process can be made to run with any required frequency by releasing the necessary energy.

Unfortunately, after analyzing for some time the essence of the work that had to be carried out *de facto* by a self-pinching beam moving along the target axis, I began to doubt the plan would succeed. The reason for my doubt was that the outlined scenario did not include something similar to a global dominant perturbation for the target nuclei “boiling” in the plasma plate. In this case, according to the logic of my own basic conception, it was difficult to rely on the appearance of conditions for the realization of a mechanism of self-organizing synthesis (presumably, it should be similar to the natural mechanism). Hence, one could not expect that the initiated process

would generate the required synthesized nuclei that were stable with respect to the action generating them and were naturally stable by possessing the maximum store of the stability to a decay.

I felt too naïve at that time to discuss my own “nuclear-synergetic projects” with my colleagues, anticipating their ironic response. At least, that’s what I thought.

All the same, at the beginning of 1998, the initiative group included the following persons: myself, heading the group; Dr. N. Tolmachev, formerly a student at the Kharkiv Aircraft Institute and then the director of a multiprofile building firm, was a sponsor of, and participant in, brainstorming sessions; active Kharkiv scientists, including an expert on nuclear physics, the owner of a huge collection of papers on a number of trends related to our interests, Dr. V. Stratienko; Professor I. Mikhailovsky; Dr. E. Bulyak, profoundly knowledgeable about beams and accelerators; Drs. V. Novikov and A. Pashchenko, the authors of numerous papers on statistical theory and thermodynamics, the theory of plasma, beams of charged particles, and nonlinear processes; Dr. I. Shapoval, an expert on mathematical modeling of physical processes and on computer structures; Kiev theorists: corresponding member of the National Academy of Sciences of Ukraine, Professor P. Fomin of the Institute of Theoretical Physics of NASU; an expert in the field of coherent processes and nuclear physics, Professor V. Vysotskii of T. Shevchenko Kiev National University, the author of one of the first models of inversionless γ -lasers. Up to the middle of 1998, the initiative group held the view that further theoretical discussions were unpromising without an experimental foundation and without the possibility to practically verify the developed ideas.

So, it became urgent to find new investors who could help in the establishment of a small research laboratory and in the creation of an experimental setup that would allow us to verify the main working hypotheses and select the viable ones from among them.

At that moment, *deus ex machina* again intervened, owing to a meeting I had with the directors of a large Kiev business concern, the Kiev Polytechnical Institute graduates Andrei Bovsunovsky and Aleksandr Kokhno, who after 1991 had left the laboratories of military plants and, together with partners, established a large-scale multiprofile holding.

After a half-year study of the problem, new potential investors were fired up by the idea and finally agreed to support our work.

We posed the following program: Establish a physical laboratory in nine months, design, produce, and launch the setup (a hard-current high-voltage generator of electric high-power pulses), provide the formation of a focused beam of electrons and, with its help, get and demonstrate the real

evidence for the attainment of the introduction of energy into a target. This last step would ensure, in particular, the fulfillment of the conditions for the positive gain of energy needed for inertial thermonuclear synthesis.

We had only nine months, and it was difficult to imagine that the allocated funds would suffice for the posed task.

However, we had no choice. Besides, I felt the inexplicable confidence, fed by a sixth sense in the saving potentiality of the general hypothesis about the principle of dynamical harmonization.

In late April 1999, due to efforts of new investors, we organized the Electrodynamics Laboratory in the structure of the Kiev company “Enran.” The purpose of the Laboratory was to realize the project which received the symbolic name *Luch*. The mission of the Laboratory was briefly formulated as follows: to create an experimental beam-based driver for inertial thermonuclear synthesis on the principles of superconcentration of the energy of an electron beam in the small internal (near-axis) volume of a thin cylindrical target.

After nine months, in January 2000, the private physical laboratory, possessing the necessary measuring and vacuum facilities, was functioning as was planned, in the leased and repaired premises of a deserted production base. We launched a generator of electric power pulses which allowed us to derive a beam of electrons with a total energy up to 300 J and a pulse duration up to 100 ns.

During this period, we carried out the initial 35 experiments — discharges with thin, up to $300\ \mu\text{m}$, target anodes.

Most members of our team believed that our goal was in sight. Very soon we would observe a thin channel along the target axis as a result of the formation of a self-pinching plasma with an ion density $>10^{24}\ \text{cm}^{-3}$ and an ion temperature $>10\ \text{keV}$. Thus, the product $n\tau$ should exceed the threshold value $10^{14}\ \text{s} \cdot \text{cm}^{-3}$ and reach a value $>5 \cdot 10^{16}\ \text{s} \cdot \text{cm}^{-3}$. What would remain was only to place a thermonuclear target on the axis, register a positive release of energy, mail the communication to *Phys. Rev. Lett.*, and confidently to await the acclaim of the scientific community.

However, the process was not running for some reason!

The beam resisted our attempt to squeeze it along the target axis and thus to create a thermonuclear plasma along the way. Moreover, we did not practically observe any evidences for the localization of a more or less significant energy in some small volume of substance. Our optimism began to wane. The experts who had recently foreseen the required behavior of a beam gave various recommendations for changes in the parameters of the driver and in the diode geometry, but then their flow of recommendations ceased and the brightness in their eyes faded.

The time given for the finding of results had passed, and the allocated funds were spent. We arrived at a dramatic collapse of our risky attempt. The “anesthetic” thought that “we are not the first, and we will not be the last” also gave no consolation.

Our investors were not indifferent observers; they asked me, as the head of the project and the Laboratory, only two questions: “What does it mean?” and “Where are your regularization and harmonization?”

What remained for me was to recognize defeat and say good-bye for ever to my beloved physics of nuclear synthesis after a fascinating but brief and unrequited fling.

However, an inner voice imposed an inexplicable calm and asserted that literally *nothing* was done to realize the idea championed by it.

I had to analyze again the reasons for our failure. To do this analysis and to make a last attempt to successfully carry out the experiment, we had two to three weeks; after the end of February 2000, work in the scope of our project had to be interrupted for a long period or for ever.

Our analysis revealed the following:

1. The localization of the focus of a superdense electron beam on the end of a target-anode is not stable. Hence, one should use a compulsory force fixation by unknown means.
2. Even if the above problem could be solved, the compression and superintense heating by the self-focusing electron beam cannot be considered a dominant perturbation common to the atoms and nuclei of a target substance, because, in this case, a coherent and unidirectional excitation of their states by a mass force is absent in principle. Moreover, the intense *heating of a substance*, only by increasing the energy of the chaotic movement of particles, *cannot play the role of a dominant external perturbation* in principle and, hence, cannot stimulate the evolutionary energy-gained fusion of the initial particles of a nuclear fuel into the more highly developed nuclear structures of synthesis products.

In other words, it was obvious that the heating of the plasma hampers the efficient and successful self-organizing synthesis of nuclei; rather than stimulating this process, it only creates the conditions for random binary nuclear collisions, only a small part of which can result in fusion. In this case, though, the reactions of synthesis for the lightest nuclei are energy-gained and are accompanied by the release of free energy; the mass defect formation *does not lead* for binary reactions to a *decrease in the inertia of the entire totality* of elements participating in the response on the external action *by any from the separated degrees of freedom in the space of states*

of the initial system of particles and hence does not correspond to the principle of dynamical harmonization.

Nothing remained to be done except for one more attempt, possibly the last, to find the “golden key” for nucleosynthesis which, on the one hand, could explain at least a part of the actually observed astrophysical phenomena related to the creation of the spectrum of the chemical elements and, on the other hand, would admit the occurrence of nucleosynthesis under laboratory conditions.

Despite the drawn-out prehistory, the fast choice of a successful, as is now clear, solution was promoted by time restrictions and the complete absence of any constructive ideas except ones not canonized in the traditional approaches to controlled thermonuclear synthesis.

It became clear that the electron beam by itself is not a coherent and monochromatic flow of energy; it transfers energy to a target for a period that is long on the nuclear scale and thus cannot play the role of a dominant external perturbation for a macroscopic ensemble of particles that could act synchronously and co-phasally on them all as a mass force.

At the same time, it is difficult to find an alternative to a weakly relativistic electron beam from the viewpoint of both the efficiency of a volume interaction with the target substance and the excitation of its collective degrees of freedom.

One day, it suddenly dawned on me, as a fully obvious thought, that the electron beam should be used for the excitation of a coherent avalanche-like self-supporting low or isentropic *secondary* (with respect to the beam) process which will develop by the laws of nonlinear phenomena with a positive feedback.

The requirements of coherence and self-preservation for the initiated secondary process imply that this process has to be wavy and soliton-like, whereas the necessity of both a continuous “sharpening” of the process and a concentration of the released energy demands that the process should be self-focusing and spherically or cylindrically (concentrically) convergent.

Intuitively, I felt the impending birth of the *conception of the artificially initiated collapse of a microtarget*, which is considered in the next chapter. This brought to a close the long prehistory of the invention of a means of shock compression of a substance, whose substantiation, experimental testing, and attempted theoretical explanation constitute the bulk of the present book.

SELF-ORGANIZING NUCLEOSYNTHESIS IN SUPERDENSE PLASMA

S. V. Adamenko

In the last decades of the 20th century, revolutionary progress has been made in studying the mechanisms of self-organization of matter (see Refs. [4–9]), using fundamental knowledge in many areas of science. Principles of self-organization developed in those studies have been successfully applied to understanding and controlling many complex processes, such as chemical reactions, laser generation, etc.

At the same time, the role of collective self-organization processes in physics of elementary particles, atomic nuclei, and natural nucleosynthesis still is not realized as being of key importance. The next years are to be marked by the ever-increasing interest in the processes of self-organization in the nuclear matter, and the change of focus from the problems related to analysis of its components towards those of finding the laws applying to the synthesis of its structures. In our view, this is the area to look for solutions to a number of fundamental physical problems.

It is well known that solutions to intricate problems are often based on fresh ideas and hypotheses that push the research in nontraditional areas. This monograph is the first presentation of the interrelated key experimental and theoretical results of the Luch project which has been no exception to the above.

Over a long period of time (since the early 1970s), researchers who later became involved in this project are gradually creating a set of working hypotheses, as well as system-level, logical, and physical models aimed at the creation of such scenarios of nuclear transformations occurring in nature that would allow the following to be done:

- to explain consistently, without adding new paradoxes while solving ones that already exist, a wider range of phenomena related to nuclear transformations observed in nature, as well as in physical experiments;
- to find such realistic approaches towards the problem of controllable nucleosynthesis that would open new ways for the creation of environmentally safe technology for the deactivation of radioactivity that

would be self-sufficient in energy terms, through a deep nuclear transformation of industrial radioactive waste, by producing a combination of stable isotopes of newly created chemical elements.

In our view, there are no fundamental obstacles preventing us from raising such a problem, since, first, it does not contradict the fundamental laws of the Nature, and, second, for the macroscopic quantity of a radioactive material, being any given mixture of isotopes (including both radioactive and stable nuclei), even without initiating its neutronization and protonization, there will always exist such a distribution of protons and neutrons (whose numbers are determined by the composition of the initial mixture of isotopes) among newly created stable nuclei that the weighted average binding energy per nucleon will be higher than that for the initial radioactive mixture, so that the redistribution will be accompanied by the energy release sufficient for its self-sustaining development.

It seems obvious that, in order to bring, in a controlled way, a macroscopic quantity of nuclei or atoms from an initial state into a final one being expedient in energy terms, one should take into account the potential mechanisms of collective nuclear and atom transformations, while the dynamical transient processes will lead to the self-organization in complex systems of nucleons or those of nucleons and electrons.

As a huge contribution to the development of the theory of self-organization in matter, there have been ideas developed by the Brussels school led by I. Prigogine (see, e.g., Refs. [4,5]). The core gist of those ideas is as follows. Nonequilibrium processes, instabilities, and fluctuations play the key role in the creation of structures in the material world, and all systems contain subsystems that keep fluctuating. Sometimes, an individual fluctuation may become so strong due to a positive feedback that the existing organization does not survive and is destroyed at a special point called the bifurcation point and reaches a higher level of the ordered organization. Prigogine has called those structures with high degree of order as “dissipative structures”.

2.1. Synthesis Process as an Instrument for Changing the Inertia of the Interactive Particles Ensemble

Around the bifurcation point, physical systems become very sensitive to even the weakest external impacts, and, in a state being far from equilibrium, such an impact may cause a rearrangement of the structure.

Note that it was usually assumed that the external controlling parameters are changing sufficiently slowly. Of course, dynamic (pulse) or stochastic changes in external controlling parameters open new

opportunities for the self-organization and application of the mathematical control theory.

The control theory is the area, where the author was also involved at the start of his scientific career. He was studying the possibilities for obtaining an analytical solution to the problem of synthesis of the optimal multilinked linear dynamical system which was analyzed over continuous or discrete time in its phase space, with the constraints set by the following equations (this consideration involves a discrete time):

$$\mathbf{P}_i^T \cdot \mathbf{x}_k \equiv 0, \quad k > k_0, \quad \text{if} \quad \mathbf{P}_i^T \cdot \mathbf{x}_{k_0} \equiv 0, \quad (2.1)$$

where $\mathbf{x}_k$ is the n -dimensional vector of phase coordinates (state vector) of the dynamical system at the k -th point of the trajectory; $(\cdot)^T$ is the transposition symbol; $\mathbf{P}_i$ is the direction vector for a hyperplane in the phase space, in which trajectories of the synthesized system should remain over the whole duration of the transient process, while the image point of the system returns to the undisturbed state or to the program trajectory, where $\mathbf{x}_k \equiv 0$.

It was proposed to use the so-called *binary synthesis method* (see Refs. 13–15).

The gist of this method is the creation of an algorithm for finding the optimal control law with feedback,

$$\mathbf{U}_{0k} = K_0 \mathbf{x}_k, \quad (2.2)$$

where K_0 is the unknown optimal control $k \times n$ matrix. This law would provide, along the trajectories of the linear dynamical system defined by the difference equations

$$\mathbf{x}_{k+1} = A\mathbf{x}_k + BK_0\mathbf{x}_k + \mathbf{q}_k\Psi_k, \quad (2.3)$$

the achievement, in the course of the transition of the system into the undisturbed state, $\mathbf{x}_k \equiv 0$, of the minimum of the *a posteriori* quality functional

$$I = \sum_{k=0}^{\infty} \left[\mathbf{x}_k^T Q \mathbf{x}_k + \mathbf{U}_{0k}^T B^T B \mathbf{U}_{0k} \right] = \sum_{k=0}^{\infty} \mathbf{x}_k^T \left[Q + K_0^T B^T B K_0 \right] \mathbf{x}_k \rightarrow \min. \quad (2.4)$$

Here, not only the matrix K_0 , but also (unlike the classical approach) matrix Q are not set *a priori* but are connected through the equation

$$Q + K_0^T B^T B K_0 = \sum_i \alpha_i \mathbf{P}_i \mathbf{P}_i^T + I, \quad (2.5)$$

with the positively defined matrix of quadric quantics of the so-called optimal goal function that reflects the optimization goals which are known

a priori, e.g.,

$$\omega_0(\mathbf{x}_k) \triangleq \mathbf{x}_k^T \left(\sum_i \alpha_i \mathbf{P}_i \mathbf{P}_i^T + I \right) \mathbf{x}_k = \mathbf{x}_k^T C \mathbf{x}_k. \quad (2.6)$$

The *a priori* goal function (Eq. 2.6) is, in its turn, nothing but a sum of independent “penalties” with the respective priorities or weight coefficients, α_i , for the deviation of trajectories of the optimal system (Eq. 2.3) from each of l set goal hyperplanes (Eq. 2.1).

In Eqs. 2.2–2.6:

- A — $(n \times n)$ matrix, the operator of the linearized dynamical system;
- B — control $(n \times k)$ matrix;
- Ψ_k — disturbance intensity at the k -th point of the trajectory;
- $\mathbf{q}_k$ — n -dimensional vector of coefficients of the sensitivity of the system phase variables to the disturbance at the k -th point of the trajectory (the disturbance direction);
- I — unitary $(n \times n)$ matrix;
- $\alpha_i > 0$ — the numerical value of the relative priority of the i -th constraint (Eq. 2.1).

The finding of the structure of optimal links for the system (Eq. 2.3) with the quality criterion Eqs. 2.4–2.6 is based on the dynamic programming methods. According to those methods (see Refs. 16–18), the optimal control law for the system (Eq. 2.3) should satisfy the conditions

$$K_0 = -(R + B^T P B)^{-1} B^T P A, \quad (2.7)$$

where, for the quality functional Eq. 2.4, $R = B^T B$, and P satisfies the matrix equation

$$\begin{aligned} P &= C + A^T P A - A^T P B (R + B^T P B)^{-1} \\ &\quad \times (I + R (R + B^T P B)^{-1}) B^T P A, \end{aligned} \quad (2.8)$$

which can be considered as a generalization of the discrete Riccati’s equation appearing in the classical variant of the optimal control problem for a linear system with quadratic functional.

Equation 2.8 can be solved using the iteration method similar to those used for the quadratic discrete Riccati’s equation, one of those methods being described in Ref. 6.

It can be proven that, if $\Psi_k = 0$, when $k > k_0$, then it is possible, through increasing the values of the weight coefficients α_i in Eq. 2.6, to

achieve the solutions to the optimal system Eqs. 2.3–2.6, whatever accuracy of the constraint (Eq. 2.1):

$$\mathbf{P}_j^T \cdot \mathbf{x}_k = 0, \quad k > k_0, \quad (2.9)$$

first of all, for such j that

$$\alpha_j \gg \alpha_i, \quad \forall i \neq j. \quad (2.10)$$

Let us consider the average disturbance with intensity Ψ and direction $\mathbf{q}$,

$$\Psi = \left\| \sum_{k=1}^{k_0} \Psi_k \cdot \mathbf{q}_k \right\|, \quad \mathbf{q} = \left(\sum_{k=1}^{k_0} \Psi_k \cdot \mathbf{q}_k \right) / \Psi. \quad (2.11)$$

The intensity of the averaged disturbance Ψ that is enough to initiate the structural changes in the system will be named as “dominating”. In Eqs. 2.1–2.9, we will assign, instead of the set of directing vectors of the destination hyperplanes $\mathbf{P}_i$, the only one vector $\mathbf{P} = \mathbf{q}$ and the single priority α proportional to Ψ . Accordingly, the matrix C in the *a priori* criterion function Eq. 2.6 will be determined by the equation $C = \alpha \mathbf{P} \mathbf{P}^T + I$.

The term “dominating disturbance” will be used below for pulse (shock) external impacts of high intensity, which disturb the elements of a complex dynamical system or its independent phase coordinates in a synchronous and cophasal (coherent) way.

The qualitative analysis of changes in the eigenvalues of the optimized system Eqs. 2.2–2.6 and the respective changes in its phase-plane portrait that occur when inequality Eq. 2.10 is strengthened (see Fig. 2.1) shows, in particular, that its proper motions relative to the set hyperplane Eq. 2.1 may have whatever small inertness, and the hyperplane Eq. 2.1 may be whatever close to the $(n - 1)$ -dimensional proper (invariant) subspace of the system, which altogether provides for the minimum value of the functional Eq. 2.4.

We note that, in the problem Eqs. 2.1–2.9, the optimal structure of the multilinked dynamical system can be found in a self-consistent way simultaneously with the parameters of its optimality criterion that are not set *a priori* and depend only on the objective factors: current parameters of the system and external disturbances that affect it.

Thus, in the framework of the binary synthesis problem¹, if we set $\mathbf{P} = \mathbf{q}$, the optimal structure can be found in a self-consistent way,

¹The binary synthesis problem was meant to be a special generalization of the classical problem of analytical design of optimal regulators (see Refs. 16–18) for those situations, where the traditional subjective selection of the matrix Q of the quadratic quality functional Eq. 2.4 (which in practice is done through many attempts) is unacceptable (we say, however, that those cases make a majority).

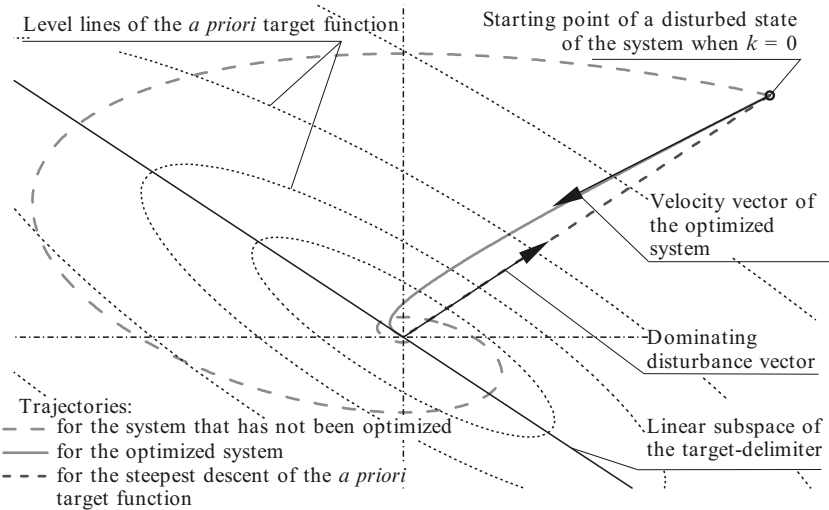


Fig. 2.1. Transformation of phase trajectories for the dynamical system as a result of its optimization using the binary synthesis algorithm.

simultaneously with finding the respective parameters (which are not fixed *a priori*) of the optimality criterion for the process Eqs. 2.4–2.10 which depends, like the system itself, only on the system's initial parameters and external disturbances that form the system. When disturbances change, then, in principle, next cycles of adaptation to external conditions can automatically start in the system. Those cycles, which repeat with every occurrence of the dominating disturbance, are able to maintain the process of continuous self-organization and reorganization in the system, providing asymptotically, in particular, the following:

- “reflection of images” of external dominating disturbances in the structure of links between the system elements and, as a result, their reflection in the set of its proper subspaces of various dimensions, which can be seen as some sort of “memory mechanism” in the system, and as the system's adaptation to dominating external disturbances
- maximization of stability and minimization of inertness of the system's forced motions that emerge in its phase space as a result of external dominating impacts

The binary synthesis algorithm will obviously have similar distinctive features of self-organization when optimizing proper motions of the system for which the following is true:

- Components of the state vector, $\mathbf{x}_k$, reflect the disturbance of the internal degrees of freedom of the dynamical system or those of the particles it comprises.
- Functional Eq. 2.4 has the dimension of energy and is the integral measure of the excitation of a system and its components, or the integral measure of the uncertainty of their state.

2.2. Main Hypotheses to the Conception of Optimal Conditions for Nuclear Synthesis

Based on the results of researches from a viewpoint of the control theory on the qualitative features of multilinked dynamical systems using methods of the control theory, as well algorithms of optimization for their phase trajectories (see Refs. 16, 17), we later analyzed the most characteristic traits of self-organization in multilinked, stable dynamical structures (systems) of any physical nature, including atomic nuclei as stable assemblies of nucleons.

This analysis resulted in understanding of the key role of the following factors in processes of such nature:

- intense *external disturbances of a certain kind*, called dominating
- general universal law, which can be called *the principle of regularization of dominating disturbances and dynamic harmonization of systems*

The dynamic harmonization principle is implemented when a set of links between source units is formed in a self-consistent way in the structure that is created or reorganized. The set of links formed under this principle minimizes or maximizes its inertness of the structure in the direction of the dominating disturbance vector in the system's phase space.

Note that, in physical terms, the problem of inertness evolution can be reduced (taking into account the equivalence between energy and mass) to the problem of evolution of the system's energy. That is, the energy of a system of particles will include, of course, the energy of their interaction, which will define the binding energy. Thus, the conclusions made using the control theory could, in principle, be applied to the problems of energy consumption and release. This idea compels us to analyze the problem of fusion from the nontraditional perspective of self-organization theory and control theory.

As a starting point for such an analysis, we take a set of basic hypotheses on the possible universal mechanisms of self-organization and reorganization in complex dynamical electron, nucleon, and nuclear structures considered from the most general perspective of the system analysis and stability theory.

The first basic hypothesis is as follows: *A set of interacting nuclei (as well as that of interacting atoms) is a dynamical system with links, which is subjected to the general laws of self-organization in multilinked, stable dynamical systems, in particular, the dynamic harmonization of systems.*

The assumption about the universal nature of the dynamic harmonization principle results in the statement that any set² of interacting elements in which old links are destroyed and/or new ones are established between elements, when external forces are applied to it, will self-consistently “determine” the optimal direction for changing its structure.

The forceful coercive creation of the subjectively “needed” structure by applying an external impact directly to the system elements or links between them (controlled synthesis) may produce the required result in the only case *where the structure being forcefully created is identical to the one that corresponds to the dynamic harmonization principle.*

Looking at the problem of the initiation of self-sustaining exoergic reactions of nucleosynthesis, we can single out what is probably the most important factor in this process: a decrease in the average and/or total mass of participating nucleons (i.e., the creation of the mass defect).

As we know, the mass of any material object (nucleon, nucleus, atom, etc.) is a measure of inertness of that object.

Following the above logic of the principle of dynamic harmonization, a solution to the problem of obtaining a negative mass defect and corresponding energy release should be found in the area of choosing an initiating mechanism (a driver) whose action would stimulate the system being reformed, which is in general an electron-nucleus or electron-nucleon megasystem—the local volume of the target source matter, precisely to decrease the average and/or total inertness (i.e., mass) of the particles affected by the dominating impact.

It is obvious that if there is no acceleration, then the motion (behavior) of the system will not depend on the inertness or masses of the particles that make up the system. Hence, a conclusion can be made that, in terms of the dynamic harmonization principle, spending the energy of an external impact (driver) for the source particles to only reach a *high final velocity or energy* would mean a failure to use the evolutionary potential of the system for its nuclear transmutation, i.e., the inefficient way of action.

The physical sense of the dynamic harmonization principle can be reduced to the following. The evolutionary synthesis using an optimal external dominating disturbance (optimal driver) is all about *initiating the proper*

²Strictly speaking, not “any” set but that whose elements have a finite inertness (which is small for the dominating disturbance) and can interact through establishing the links of some physical nature and changing the intensity of those links.

motion of the ensemble of interacting source elements (particles) into a new state in terms of both energy and topology, and, accordingly, to a new required structure, whose motion will be along the steepest descent path, with minimization of energy or another more general functional spent during the transition from the initial to the final state.

In the general case, we are talking about a new approach towards the synthesis of multiparticle multilinked stable systems or structures of any physical nature. The distinctive feature of this new approach is the thesis that, in the Nature, the synthesis of a system or the self-organization of its structure is always a self-consistent collective response of elements in the source ensemble to the common dominating force which disturbs their state of rest or a nonaccelerated motion in the system's eigenfunction space. The common dominating disturbance coherently transfers the momentum p_D of the unidirectional motion in the phase space to all elements of the ensemble, the momentum being satisfied the condition

$$|p_D| \gg |p_{T\max}|, \quad (2.12)$$

where $|p_{T\max}|$ is the maximum absolute value of the momentum for the proper (uncoordinated) thermal motion of any element of the ensemble in the same phase space over the whole period of transition from the initial state (before the common dominating disturbance appeared) into the steady final state.

A fundamental distinctive feature of the proposed new approach toward the synthesis is the search for and the selection of such a mechanism of initiating the self-sustaining and self-consistent process using the criterion of optimality of the *process of transition of the dynamic system (ensemble of interacting particles)* from the initial state/configuration into a new state/configuration *being "energy-efficient" not only at the final point, but also integrally over the whole duration of the transient process.*

The second basic hypothesis is as follows: *With increasing the number of the elements being excited in the synthesized system, there is a fast growth in the influence of cross-links of small intensity, the establishment of such links resulting in decreasing the inertness of the system in the designated direction, i.e., that of the impact direction.*

It is possible to increase the number of interacting particles of the matter to any value by transforming the matter into a critical state, where the particles' behavior is as if each of them "feels" the current state of all other particles (see Ref. 6), i.e., the characteristic correlation length under the phase transition in the system tends to infinity. The physical nature of this generally recognized "feeling" has not been studied completely as yet; this fact, however, does not prevent us in any way from taking this effect

into account when arranging the initiation of a self-organizing process of nucleosynthesis.

We have assumed that the effective range of interaction for nuclear forces is determined by the degree of coherence of dominating disturbances that initiate the interaction between particles. Particles in the matter are able to interact, including the exchange interaction and, respectively, the collective behavior in any of the possible critical states of the matter. The fundamental importance will be attached to the right definition of which critical state of the matter is optimal (or at least favorable) for the initiating and the intense running of self-sustaining collective nuclear reactions.

States on the liquid–gas or gas–plasma interface are indeed critical; however, in those states, “criticality” is applied to electrons only, as nuclei, while they still keep electron shells, are not in a critical state, because each of those nuclei stays in its own “rest niche”. The nearest critical state where those niches are common is the electron-nucleus plasma, i.e., the state, where the ionization number for atoms is equal to their nuclear charge.

In this case, because of the dynamic harmonization principle, nuclei involved in the collective are “faced with the problem” of choosing the “way towards a new stability”, the options here being (1) back to the formation of nonionized atoms or (2) forward to the creation of new nucleon ensembles or new nuclei.

It is important to note that, in the state of electron-nuclei plasma, electrons are, like nuclei, elements within the ensemble that has to “make its choice”. Therefore, if their participation in the collective nuclear process may contribute to its efficiency, i.e., the creation of the maximum mass defect within the minimum period of time, then they will participate in the process of collective interaction between particles in the system, even without external impact, under the self-organization, or evolutionary synthesis, scenario.

The third basic hypothesis is as follows: *The required synergetic process of collective transmutation of nuclei is nothing but the natural nuclear combustion of a collapsing matter. Its essential features are identical to those of the natural nucleosynthesis (where the whole range of the stable isotopes of chemical elements is created). This is a fundamentally multiparticle self-organizing process which cannot be implemented through pair collisions between particles, and thus it is fundamentally impossible to reproduce it in any conditions other than the collapse (the compression to extreme densities) of macroscopic quantities of matter.*

When applied to explosive astrophysical processes, such as the collapse in a burned massive star resulting in a supernova burst, the hypothesis

has allowed us to assume that the stable isotopes of chemical elements found in the Nature are not created through a long decay of their radioactive predecessors, which, in turn, have been created in the Big Bang. Instead, those stable isotopes themselves are immediate products of the natural nucleosynthesis which develops in the following ways:

- either as a collective, exoergic process of *nuclear combustion of matter* in a collapsing supernova
- or, as a process of decay of superheavy nuclei (superdense matter) created as a result of the collapse in massive stars via the cluster radioactivity mechanism

There is no need to spend much time looking for candidates to become a *natural* driver for nuclear combustion, since such a driver is obvious. It is the universal gravitation which mutually attracts huge amounts of matter and results in the development of a gravitational collapse.

It does not seem possible to use this kind of a driver under the earth conditions, but there is no fundamental reason to expect that the same final outcome cannot be achieved using a different tool, i.e., a different driver. It is clear that in order to do that, we need to find some key distinctive feature of the natural collapse mechanism we are trying to substitute, and then this distinctive feature has to be reproduced, without reproducing the whole mechanism.

As a result of analyzing those peculiarities of gravitational forces initiating a collapse in massive space objects, which are significant in terms of the efficiency of the required process, the fourth basic hypothesis has emerged. It consists of the following two interrelated statements:

1. *The most important property of the gravitational collapse is its being initiated by a self-generated common dominating disturbance which has character of a mass force and emerges due to the coherent amplification of gravitation effects for a macroscopic ensemble of particles of matter located at the same distance from the center of mass of the collapsing object, i.e., those located in the volume of a thin spherical layer (see Ref. 19);*
2. *Collective self-consistent coherently accelerated motion of a set of particles of matter, where the centripetal component of the momentum of each particle is much higher than its thermal (chaotic) component, brings a portion of matter made of those particles into a special “collectivized” critical state.*

In Part IV, we will show that it can be formally drawn from solutions to the kinetic equations for the particle system under a dominating

impact. Solutions to a number of approximate forms of kinetic equations (Boltzmann, Landau, Lennard–Balescu, and nonlinear Fokker–Planck equation) under the condition of a constant flux in the phase space (the existence of a dominating impact in the system!) have power asymptotics, i.e., states characterized by strong large-scale correlations.

In such a state with strong correlations, the optimal conditions may emerge for the avalanche-like spontaneous creation of highly organized multiparticle multilinked dynamical network structures. In thin macroscopic structured spherical shells with a common center of mass, the dynamical structures may emerge on the basis of all kinds of physical interactions (including the strong and Coulomb ones), for which the necessary conditions exist (within the respective shell where the matter is in the critical state). Those processes should probably depend on the reaching of the threshold densities by the matter within the shell in the state of electron-nucleon-nucleus plasma.

In this case, because the energy of the collective interaction grows with the number N of interacting particles, *the energy of the “global” interaction of particles in a self-structurizing shell at sufficiently great N can significantly exceed the energy of their “local” interaction and can define the topology of a spontaneously self-organizing branched macroscopic nuclear structure which minimizes or maximizes the own total inertia relative to the action of a centripetal mass force (i.e., which minimizes or maximizes the own inertial mass) and, respectively, maximizes or minimizes the arising mass defect and the energy of a global coupling of all particles of the shell.*

Based on this hypothesis, we have finally come to an assumption that in the process of gravitational collapse in a massive burned star, the decisive role for the self-sustaining process of exoergic nuclear combustion may be played in some cases by a *solitary spherical wave of the matter (energy) density collapsing to the center*, such a wave being formed on the surface of the collapsing body through one physical mechanism or another.

As a result of the nonlinear increase in the steepness of the leading edge of this collapsing wave and the increase in the matter density in the wave to the critical point, there is a possibility for the spontaneous start of the collective processes of formation of a macroscopic electron-nucleus structure. In that structure, there are coherent states of interacting particles along the radial coordinate $r(t)$, their states being at the same time *strongly correlated* (see Ref. 20) in the 2D subspace of angular coordinates ψ , φ (the shell surface subspace).

In this case of the concerted collective behavior of quantum mechanical particles in the volume of a collapsing spherical shell layer, we can use

the uncertainty relations for strongly correlated states (see Refs. 20, 21):

$$\begin{aligned}\sqrt{1-k^2} [\Delta\psi \cdot r(t)] \cdot \Delta p_\theta(t) &\geq \hbar/2, \\ \sqrt{1-k^2} [\Delta\varphi \cdot r(t)] \cdot \Delta p_\varphi(t) &\geq \hbar/2,\end{aligned}\tag{2.13}$$

where $0 \leq k^2 < 1$ is the correlation coefficient, $(\Delta\psi \cdot r)$, $(\Delta\varphi \cdot r)$, and Δp_ψ , Δp_φ are the variances of linear shifts over the angular coordinates and those of the respective momentum values of an arbitrary particle that participates in the transfer of part of the object's mass towards its center by the collapsing wave.

It is obvious that the states of particles in this shell system become more correlated as the system further deviates from the equilibrium, energy flows increase in the phase space, thickness of the collapsing shell decreases, and density of a substance in its volume increases with time, i.e., $k(t) \xrightarrow{t \rightarrow T} 1$.

According to the energy conservation law, Eq. 2.13 suggests that simultaneously:

- the variances of the angular coordinates $\Delta\psi$, $\Delta\varphi$ unlimitedly increase
- the range of action for nuclear forces in the subspace ψ , φ sharply increases due to the “smearing” of the range of definition of a wave function for each of the particles in the strongly correlated space over the 2D surface of the shell

As a result, the conditions emerge that are necessary and, due to the effect of the dominating disturbance, also sufficient for the *collective nuclear interaction* between particles in the shell, which is accompanied by a mass defect and a further increase of the matter and energy density in the shell. The positive feedback by the density, which determines the behavior of the collapsing shell, strongly increases.

However, taking into account that $\Delta\psi \cdot r(t) \leq 2\pi r(t)$ and $\Delta\varphi \cdot r(t) \leq 2\pi r(t)$, due to Eq. 2.13, when $t \rightarrow T$ and $r(t) \rightarrow 0$, the variance of the particles' coordinates in the 2D subspace of the shell surface is forced to decrease rapidly; in the situation where the correlation between particle states in the shell strengthens, this must result, due to the obvious relation

$$\Delta p_{\psi,\varphi}(t) \geq \frac{\hbar}{4\pi r(t)\sqrt{1-k^2}},\tag{2.14}$$

in increasing the variance and average values of their momenta and the increase in the kinetic energy.

Source of the free energy that is released is the emergence of a negative integral mass defect in products of the collective nuclear transformations in

the volume of the collapsing wave-shell and in the volume of the collapsing object that has been scanned by the wave. The integral momentum of the particle system is conserved, since the wave-shell collapses concentrically and the velocity of its center of mass does not change.

On the other hand, if the particles in the ensemble are in a coherent state in the 1D subspace $r(t)$, then the uncertainty relation between the coordinate and momentum is known to become the equality

$$\Delta r(t) \cdot \Delta p_r(t) = \hbar/2. \quad (2.15)$$

From this equality, it is even more obvious that since $\Delta r(t) \ll r(t)$, the variance of the momentum $\Delta p_r(t)$ will quickly increase with decrease in $r(t)$. As $r(t) \rightarrow 0$, it can become whatever large, which will be unavoidably accompanied with a fast increase in the kinetic energy of motion directed towards the center of the collapse for particles that currently make up the shell.

The sources of the released energy here probably include, like in the previous case, on the one hand, a decrease in the potential energy of particles of the shell matter, and, on the other hand, a fast decrease in the free surface area with a respective decrease in the surface energy of the nuclear megastructure in the shell volume.

Since the evolution of a wave-shell which collapses to the target center is assumed to be a stable self-organizing process which is always completed in a small vicinity of the own focal point (the singular point), we are compelled to make attempt to answer the following important questions:

- Why does not the release of the nuclear combustion energy in the shell volume lead to its overheating and the fracture on the early stages of the development of the process?
- Where and how is the nuclear combustion energy accumulated during the entire time of the evolution of a wave-shell from the time moment of its transition to the active phase of nuclear transformations occurring in its volume to the time moment of its collapse at the center?

The same questions can be formulated in a more specific and, simultaneously, more general form: How can the processes of release and absorption of energy coexist in the frame of one megacuster electron-nuclear system and what can define the energy directivity of the collective nuclear transformations in a superdense substance when the problems of the Coulomb repulsion of nuclei and the smallness of the action radius of nuclear forces become insignificant?

By following the logic of our conception, we are compelled to seek the answer to these questions with the help of the principle of dynamic

harmonization and its simplest model used in linear problems of the binary synthesis (2.2)–(2.6).

This principle yields that the formation of the ensembles (clusters, groups, etc.) of elements, particles, and nucleons in a self-organizing set should be subordinated to the problem of minimization of functional (2.4) on the trajectories of a perturbed (excited) motion of system (2.3).

To reduce the volume and to enhance the transparency of the forthcoming interpretation, we use a slightly transformed continuous representation of system (2.2)–(2.6), both being quite equivalent in the context under consideration:

$$U_0(t) = K_0 X(t), \quad (2.16)$$

$$\dot{X}(t) = F \{X(t), \Psi(t), U_0[X(t')]\}, \quad (2.17)$$

$$I = \int_{t_1}^{t_2} T[X(t)] dt \rightarrow \min. \quad (2.18)$$

Here, $T()$ is the square form of phase coordinates which represents the sum of the kinetic energies of the excited (nonequilibrium) motion of all particles of the system. These particles participate in the nuclear transformation and change, by virtue of this, their own inertia and the inertia of the collapsing shell as a whole at the expense of their mass defects:

$$T(X) \equiv \sum_{i=1}^N \frac{m_i(t) v_i^2(t)}{2}, \quad (2.19)$$

$m_i(t) \equiv m_{i0} + \Delta m_i(t)$ is the variable rest mass of the i -th particle which belongs to the composition and the structure of the shell at the time moment t ;

$\Delta m_i(t) = f\{U_0(X)\}$ is the current rest mass defect of the i -th particle which participates in the collective process of transfer of the substance-energy of the collapsing shell;

$U_0(X)$ is the operator of cross-couplings (interaction) between particles which define the current, optimal in the sense of (2.18), electron-nucleus structure of the collapsing shell and its components and, hence, the character and the energy directivity of collective nuclear transformations in it.

In view of the definition of a dominant perturbation, which possesses the character of a short-term mass force (impact), being constant on the time interval of its action on each particle separately, we may consider that a particle, through which the density wave passes, shifts along the shell radius and makes transition in the time interval $(t_2 - t_1) = \Delta t$ from the rest state at the point A (prior to the passage of the wave) in a new rest state at the point B (after the passage of the wave), by successively changing three phases of the motion:

- the phase of collective impact acceleration in the interval $(t_1, t']$ upon the passage of the leading edge of the wave
- the phase of stationary nonaccelerated motion on the wave crest (interval $[t', t'']$)
- the phase of collective impact deceleration in the interval $[t'', t_2)$ upon the passage of the trailing edge of the wave

As a result of the calculations which are executed by us³ and are omitted because of their absolutely clear logic and regularity, we may assert the following:

- Functional (2.18), being (by virtue of the principle of dynamic harmonization) the action integral only for excited motions, is transformed into the sum $(I_1 + I_2)$ of two functionals obtained for the intervals $(t_1, t']$ and $[t'', t_2)$, respectively.
- A minimum of the functional I_1 is attained by the maximization of the inertia of products of the nuclear transformation in the interval $(t_1, t']$, i.e., the nuclear transformation of a substance on the leading edge of the wave must be accompanied by the formation of a positive rest mass defect of products of this transformation $\Delta m_i > 0$, hence, by the energy absorption (endoergic process).
- In the interval of the stationary motion $[t', t'']$, the integrand of functional (2.18) (the kinetic energy of the excited motion) is equal to zero by definition; the rest mass defect of nuclear structures is not changed ($\Delta m_i(t) \equiv 0$), the accumulation of products of the nuclear transformation “coming” from the leading edge of the collapsing wave can occur.
- A minimum of the functional I_2 is attained by the minimization of the inertia of products of the nuclear transformation in the interval $[t'', t_2)$, i.e., the formation of a negative rest mass defect of particles of the substance (atoms), being the products of the collective synergetic regeneration, must happen on the trailing edge of the collapsing wave,

³With the independent help of Professor V.I. Vysotsky and Dr. V.E. Novikov.

$\Delta m_i < 0$. Hence, the energy is released (the collective exoergic process is the nuclear combustion of a substance coming from the volume of the wave-shell crest).

2.3. About the Possible Scenario of the Self-Organizing Nucleosynthesis in the Collapsing Wave of Nuclear Combustion

Using the above basic hypotheses, we can state a scenario of the process which has its internal logic, while being paradoxical in that its space-time physical model seems to be rather obvious, while the respective mathematical description, if created, is unlikely to be adequately transparent as compared with the simplicity of the main idea.

An attempt at summarizing the key points of the above basic hypotheses results in the following generalizing assumption. *The synergetic process of collective transmutation of nuclei* occurs during the natural *nuclear combustion of matter*, which is self-created and self-sustained in the evolving collapse of a macroscopic spherical wave-shell with the extreme density of energy and matter. In the volume of that wave-shell, interacting particles make up a multilinked *dynamical system*, which is self-organizing in response to the spontaneous emergence and effect of the self-reproducing *common dominating impact* that is sharpening in a nonlinear way and creating the required conditions for self-sustaining nuclear combustion and accumulation of the released energy:

1. Transition of a macroscopic portion of matter into a “collectivized” *critical state*, which is characterized by the *coherence* between the currently interacting particles along the radial coordinate (in the radial subspace), and a strong correlation between the states of the same particles in the 2D subspace of angular coordinates of the surface of the shrinking (collapsing) shell;
2. Excitation of a collective exoergic detonative cumulative chain electron-nucleus process in the body of the collapsing shell, through which the system of links between the initial elementary components of the matter would be reorganized, and the basic eigenstates of these components would change (including the proton and neutron state of nucleons). That reorganization develops towards the optimization of the self-organizing transient process of “reflection” of the common dominating disturbance.

In this case, the disturbance itself reproduces successively the excitation of a macroscopic ensemble of strongly interacting and strongly correlated particles in the scanned monolayers of the target matter. When these particles pass through the collapsing wave-shell scanning

them, they turn out first (in the passage through the leading edge of the density wave) under conditions of the collective coherent acceleration and then (in the passage through the trailing edge of the wave) under conditions of the collective coherent deceleration. This leads to a reorganization of links between the elementary components of the matter and a transformation of nuclei, which results on the one hand in

- the energy release (the nuclear combustion of a substance) on the trailing edge of a density wave and the accumulation of exoergic products of this process in a part of the body's volume already scanned by the wave, i.e., the object of a collapse;

and on the other hand in

- *the accumulation of the released energy* in endoergic products of the nuclear transformation which form and supplement the mass of a substance, being in the stationary motion and forming the crest of a collapsing density wave. This wave moves with minimum dissipation of energy to the symmetry center (the singular point) of the process during the entire course of its evolution *on the implosion stage*.

Let us try to construct a scenario of that process which would satisfy the requirements of the above basic hypotheses. We will state that scenario in the form of a step-by-step sequence of key events that develop during the hypothetical process of artificially initiated nuclear combustion.

Step 1. Using some low-inertia coherent driver, a high-power symmetrized shock impact is applied to all particles of the surface layer of a centrally or axially symmetric energy-concentrating target.

Step 2. Due to the driver impact, the substance density leaps in a thin near-surface layer, and, within that layer, particles in the substance acquire the momentum directed towards the center of the target. A closed solitary nonlinear wave emerges which is collapsing both in its thickness and its radius.

Step 3. The energy of the impact, which is transferred by the wave, is quickly redistributed over the radial coordinate of the wave-shell, a steep front edge being created with extreme gradients of both the electron and ion density.

There appears a closed self-organized layer of disturbed particles that is moving towards the target center. On its inner surface, the concentric monolayers of matter undergo successively the shock acceleration. In terms of the multilinked dynamical system which obeys the dynamic harmonization principle and can be optimized, that process means the formation of a common dominating disturbance that recurs and replicates many times. The

response of excited particles to that disturbance, i.e., transition in which the particles and their phase coordinates move towards a new stable state, can be described in the simplest case with the following equation:

$$X(n+1) = F\{X(n), K(n)X(n)\}, \quad (2.20)$$

where $X(n) = \{x_1(n), x_2(n), \dots, x_L(n)\}$ is the R -dimensional vector $\left(R \triangleq \sum_i^L M_i\right)$ of the set of excited states (excitations) of L particles; x_i is the M_i -dimensional vector of excitations (excited phase coordinates) of the i -th particle; $K(n)$ is a skew-symmetric $(R \times R)$ -matrix of links between interacting items (excitation exchange matrix), where the elements located symmetrically around its main diagonal are of the same absolute value but opposite signs, thus ensuring the conservation of the total energy of excitation transferred to the system by the external dominating disturbance; n is the discrete time; F is a function that reflects the wave nature of the process, in which the energy/mass is transferred by the collapsing layer of excited superdense matter, compliance with conservation laws, as well as the continuity and direction of the process of sequential shift of target matter monolayers towards the target center, following the collapsing wave-shell.

Step 4. A self-organizing and self-sustaining process is initiated, in which the optimal structure of links (matrix of interactions) is formed. That matrix minimizes the functional Eq. 2.4 or 2.18 that has the form of the excitation integral, and depends on the parameters characterizing the instantaneous (current) inertness of the system Eq. 2.20 or, in other words, on the roots of the system's characteristic equation, those roots being defined by the type of the function F and that of the matrix (the operator) $K^\circ(n)$.

The analysis of the dynamical system Eqs. 2.3, 2.4, or 2.17, 2.18 allows a conclusion that the efficient minimization of the excitation integral (actional integral) within the duration of the process evolution can be ensured through the establishment of such links between the particles involved in the energy transfer that would simultaneously:

- minimize the inertness of the shell
- maximize the total mass of the substance transferred by the wave, by increasing the number of particles in that wave

Thus, the principle of dynamic harmonization can be interpreted as the least-action principle or as the Prigogine principle of the minimum production of entropy concerning the problem of optimization of the inertia or the optimization of a mass defect in a dynamic system formed by the collection of inertial elements with optimizable couplings between them. The regular following to the principle of dynamic harmonization can ensure the

self-sustaining character of the nuclear transformation in a stably collapsing wave with the extremal density of a substance which simultaneously realizes in the process of its evolution:

- exoergic regeneration (nuclear combustion) of the substance of a body-object scanned by it
- transformation of the released energy of the nuclear combustion in the internal and kinetic energies of products of the endoergic nuclear transformation which form the wave crest, are accumulated in it, and form the developed macroscopic nuclear structure of a collapsing shell self-organizing, in particular, also by the criterion of maximum of the stored energy transferred to the center (to the focal point)

The final (explosive) stage of this process can also be analyzed using the above approach. Such an analysis shows that the dissipation of the energy released in the target center when the structure of the collapsed shell is destroyed can be ensured through the following:

- maximum absorption of the free energy, which has been accumulated by the wave-shell, in the nuclear transformation of its matter, when the nuclei created in the endoergic (endothermic) process of the shell fragmentation, which is inverse to the explosion, have lower binding energy than those of the initial substance of the target
- carrying of the excessive energy away out of the system by particles which weakly interact with the substance
- emergence of a low-entropy (isoentropic) wave with the energy and matter transfer from the collapse of the “hot dot” (HD) to the surface of the exploding target and then to the surrounding space

In our opinion, taking into account all those *a priori* assumptions (basic hypotheses) limits significantly the variety of possible drivers that could be used to activate simultaneously, all the physical and synergetic mechanisms considered in those assumptions in the experiment.

Let us summarize again the key requirements for such a driver:

- ability to quickly reach the high volume density of the power input in the surface layer of the target (*shock action*)
- synchronous and cophasal impact on a large ensemble of particles ($N > 10^{12}$) (*coherence*)
- ability to reach the spherical or cylindrical symmetry of shock compression of the target material (*concentricity*)

In order to estimate the degree of conformance of an external impact to the notion of common dominating disturbance, we need to introduce some

quantitative criterion of its shock nature and coherence. We will try and construct the first approximation for such a criterion based on the general considerations being in accordance with the above basic hypotheses of our concept.

Taking into account the need to provide a pulse impact on the target surface, let us consider the energy flow with its intensity sharply increasing in time and only within the time period Δt_p where the beam power increases.

The heuristic criterion to reflect simultaneously the strength, locality, and coherence of the shock impact can be expressed as

$$J \triangleq K_1 \cdot K_2 \cdot K_3, \quad (2.21)$$

where K_1 , K_2 , and K_3 characterize, respectively, the intensity, locality, and coherence (synchronism) of the shock impact. Taking into account the simplifying assumptions that are acceptable to estimate J , they can be represented by the following values:

the average speed K_1 of a change in the energy flow within the time interval where it increases

$$K_1 \triangleq \Delta P \cdot \Delta t_p^{-1}; \quad (2.22)$$

the value of K_2 inversely proportional to the part of the target volume which absorbs the shock energy:

$$K_2 \triangleq (S \cdot v_M \cdot \Delta t_p)^{-1}, \quad (2.23)$$

where S is the surface area that accepts the shock energy, v_M is the excitation propagation speed in the target body;

and K_3 inversely proportional to the shock impact duration:

$$K_3 \triangleq \Delta t_p^{-1}. \quad (2.24)$$

From Eqs. 2.21–2.24, we obtain

$$J \triangleq (S \cdot v_M)^{-1} \cdot \Delta P \cdot \Delta t_p^{-3}. \quad (2.25)$$

J has the dimension of $\text{W} \cdot \text{cm}^{-3} \cdot \text{s}^{-2}$ and characterizes the acceleration of the volume density of power of the shock impact absorbed by the surface layer of the target matter.

Under our basic hypotheses, in order to trigger a self-sustaining mechanism of dynamic harmonization in the ensemble of interacting particles, it is required that the energy of the external dominating disturbance (massive dominating force) be applied to the particles that constitute a system

(portion of matter) in a critical state. This, in its turn, needs the impact energy to exceed the ionization energy for the respective portion of the target matter, i.e., the following condition should be met (under the simplifying assumption of a linear increase in the shock beam power):

$$0.5P_{\max}\Delta t_p \geq E_i n S v_M \Delta t_p, \quad (2.26)$$

where E_i , in eV, is the specific energy (per atom) of ionization by compression which is defined by the conditions under which the process is triggered; n , in cm^{-3} is the volume density of the target material; and $V_M \Delta t_p$ is a value measured in nuclear diameters $D_n \approx 1 \text{ fm}$, and it should exceed the required multiplication factor K (which is set *a priori*) of matter density increase in the wave-shell when its radius and thickness, which are decreasing, reach their threshold values set *a priori*. The ratio Eq. 2.26 then can be transformed into the form

$$0.5P_{\max}\Delta t_p \geq E_i n S D_n K. \quad (2.27)$$

Since there is also the requirement for the matter density to reach, as a result of the nonlinear increase in the wave-shell steepness in the target, the extreme density on the leading edge of the wave, the following obvious condition should be met:

$$v_M \Delta t_p \ll R_g, \quad (2.28)$$

where R_g is the calculated radius of the nuclear combustion area in the focal volume of the target.

As a result, we obtain the following conditions required to initiate a self-developing collapse of the wave-shell when the energy of the directional shock impact is transferred to a thin surface volume of the target matter:

$$(S \cdot v_M)^{-1} \cdot \Delta P \cdot \Delta t_p^{-3} > J_{cr} \quad (2.29)$$

$$2 \frac{E_I n_i S D_n K}{P_{\max}} \leq \Delta t_p \ll R_g V_M^{-1}, \quad (2.30)$$

with the following natural conditions:

$$D_n K \ll R_{tg}, \quad (2.31)$$

$$R_g < R_{tg}, \quad (2.32)$$

where $\Delta P \approx P_{\max}$ is the increment in the flow of particles within the impact duration; J_{cr} is the lower critical threshold of efficiency (coherence) of the shock impact for a given target material.

The above basic hypotheses and required conditions Eqs. 2.29–2.32 are the basis for the development of a method for shock compression of matter, as well as the creation of a respective device (driver) for a practical implementation of the conception of artificially initiated collapse and controlled nuclear synthesis. The experimental checking done in 2000–2004 has mostly confirmed the validity of the basic hypotheses and enabled us to develop, based on experimental data, the following generalized idea of a macroscopic scenario of detonative nuclear combustion of matter in a collapsing wave of superdense plasma.

Stage 1. Emergence and hydrodynamic evolution of a collapsing wave-shell of superdense plasma

- Short-term impact of an electron beam on the near-surface layer of a solid spherical target.
- Ionization and pulsed coherent compression of a thin spherical layer of the target in the radial direction.
- Formation of a spherical soliton-like nonlinear wave (wave-shell) of the dense plasma ($\rho_n > 10^{23} \text{ cm}^{-3}$) and start of its centripetal motion that accelerates in a self-consistent way (shrinking, collapse). That motion is accompanied by the avalanche-like steepening of the front edge (decrease of its spatial length), and the quasi-isoentropic increase in the density of energy and matter, as well as in the degree of ionization at the leading edge of the wave.
- Development of an instability in plasma waves (while a nonlinear wave of density exists) and the emergence of a double-layer-like plasma-field structure in the collapsing wave volume with extreme radial field intensity in this self-organizing spherical “virtual capacitor”.
- Criticality reached by the following parameters of matter both in the wave volume and on its leading edge:
 - extreme density
 - full ionization (its degree to become equal to Z)
 - pulsed coherent superacceleration of particles in the next spherical layer that is involved in the interaction with the steep leading edge of the density wave that runs onto that layer
- The transition of the macroscopic ensemble of particles of the wave in a strongly correlated state in the subspace of the angular spherical coordinates of the incident shell and, simultaneously, in a quasicohherent

state in the radial direction, i.e., in the wave propagation direction; the appearance of the initial conditions for a collective nuclear regeneration of a substance in the collapsing density wave.

***Stage 2.** Self-sustaining pycno-nuclear combustion and the accumulation of energy in the nuclear structures of a wave with extreme density*

- Detonative nuclear self-ignition of matter in the successive spherical monolayer that undergoes a shock acceleration: the process takes the form of active pycno-nuclear combustion (the self-consistent mode of collective exoergic nuclear reactions of synthesis-fission in superdense electron-nuclei plasma, those reactions being accompanied by the creation of the integral negative mass defect in the system of nuclei and particles produced by combustion).
- Spontaneous creation of a ramified structure of a macroscopic electron-nucleon-nucleus megacluster in the collapsing wave-shell. Active process of detonative nuclear combustion of matter in the collapsing wave. Self-sustaining increase in the number of particles, as well as in integral and specific energy of the collapsing macroscopic cluster.
- The transition of a wave-transported substance in the critical coherently correlated state of the electron-nucleus plasma near the density maximum on the leading and trailing edges of the wave, which ensures the deformation (dilatation) of the domains of definition of the wave functions of particles in the subspace of angular coordinates (on the wave-formed surface of a thin-walled shell with extreme densities of substance and energy), and the relevant increase in the action radius of nuclear forces in the 2D subspace of angular coordinates against the background of a simultaneous decrease in this radius in the radial direction up to the physical minimum defined by relation (2.15).
- Self-consistent initiation, development, and running of collective nuclear transformations with opposite energy directivity: at the trailing edge of the wave in the region of intense coherent deceleration of particles, in the direction of a decrease in the integral and specific inertia (the rest mass) of formed nuclei, with the formation of a negative mass defect and the release of free energy; at the leading edge of the wave in the region of intense coherent acceleration of particles, in the direction of an increase in the integral and specific inertia (the rest mass) of products of the nuclear regeneration, with the formation of a positive mass defect and the absorption of free energy.
- Formation of the balance of the energy released on the trailing edge and absorbed on the leading edge as of the current time moment with the use of their difference for the acceleration or deceleration of the centripetal motion of the wave.

- Accumulation of the free energy, which was released in the process of nuclear combustion on the trailing edge of the wave, in the nuclear substance and in the nuclear structures of its leading edge at the expense of the processes of neutronization, the formation of nuclear structures with minimum specific binding energy per nucleon, and the increase in the Fermi energy of the degenerate electron gas.
- Stabilization and *self-consistent optimization (dynamic harmonization)* of the processes of nuclear combustion and self-development in the cluster in accordance with the following *criterion*:

$$\Delta^m N_{pr} \cdot \overline{\text{SBE}}_{pr} - \Delta^m N_{tg} \cdot \overline{\text{SBE}}_{tg} + (N_w + \Delta^m N_w) \Delta^m \text{SBE}_w \rightarrow \max, \quad (2.33)$$

where

$\Delta^m N_{pr}$ is the total number of nucleons in the nuclei produced by the pycno-nuclear combustion, which are evaporated by the trailing edge of the wave-shell at the m -th elementary step of collapse. Each elementary step means that the leading edge of the wave moves by one interatomic distance;

$\Delta^m N_{tg}$ is the total number of nucleons in the successive monoatomic layer of the target that is involved in the wave at the m -th elementary step;

$\Delta^m N_w$ is increase in the number of nucleons in the megacuster of the collapsing wave at the m -th elementary step;

$\overline{\text{SBE}}_{pr}, \overline{\text{SBE}}_{tg}$ are respectively, the weighted average binding energy per nucleon in the nuclei evaporated by the trailing edge of the wave and in those consumed by its leading edge at the m -th elementary step;

$\Delta^m \text{SBE}_w$ is increase in the specific binding energy in the megaatomic cluster of the wave-shell at the m -th elementary step of the collapse.

As an additional criterion of optimality for the sustainable self-development of the macroscopic cluster, there is the weighted average *specific binding energy* per neutron in the nuclides evaporated by the trailing edge of the shell.

- End of the active phase of the *implosive stage* of the collapse; the energy accumulated in the volume of the pycno-nuclear wave-shell

reaches its relative physical limit which is determined by the following factors:

- state of matter
 - current kinetic and potential energy of the cluster
 - characteristic time of nuclear transformations
 - current degree of correlation and coherence for the state of particles in the thin collapsing spherical (or cylindrical) layer
 - weighted average specific binding energy for source nuclei in the collapsing target
- Degeneration of the shell (the regeneration of the shell itself into a hollow sphere, in which the hollow radius at the center decreases much more rapidly than the sphere radius); the breaking of the earlier conserved balance between the release of energy (at the trailing edge) and its absorption (on the leading edge) in favor of the release due to the degeneration of the shell; the accumulation of the inabsorbable energy in the shell's body (the heating of the shell); the explosive attainment of the relative maxima of the density and the degree of protonization of a collapsed substance in the vicinity of the focal point of the collapse; the transformation of the energy accumulated in the wave into the extreme density and the excitation energy of a supercompressed macroscopic cluster-fireball; the transition of the process into the explosion stage.
 - Fragmentation and explosion of the overexcited fireball.

Stage 3. Fragmentation of the collapsed megacuster

- Formation of the mass and charge spectrum for the particles that are fragments of the exploding megacuster-fireball as a function of energy density and neutron concentration in it (Fig. 2.2).
- Formation of the final energy balance for the implosive stage of the process as a function of the whole spectrum of nuclei (from the lightest to superheavy ones) produced by the fireball explosion and those evaporated on the trailing edge of the shell over the whole duration of its evolution during the implosive stage.

Depending on the whole spectrum of nuclei born by the fireball, there may be a qualitative difference in a further evolution of the process, in terms of the released free energy (i.e., the energy not bound by created “energy-consuming” nuclei).

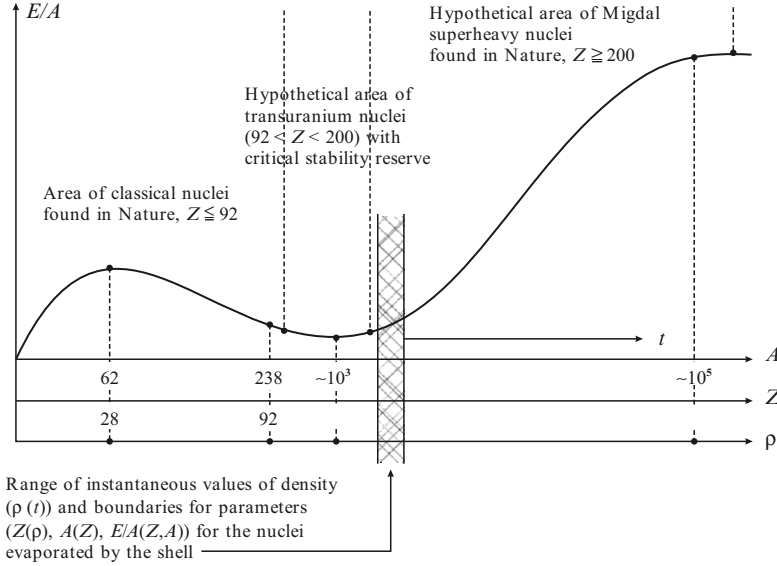


Fig. 2.2. Estimated relationship between the maximum values of parameters ($Z, A, E/A$) for nuclei evaporated by the shell and the current density of matter in the wave-shell volume.

Under the critical scenario of self-developing collapse, the integral energy release in the process may be close to zero, i.e., the following relation would be satisfied:

$$\sum_{i=1}^K \Delta \bar{E}_i \cdot N_i = \sum_{j=1}^L \Delta \bar{E}_j \cdot N_j, \quad (2.34)$$

with the condition

$$\sum_{i=1}^K N_i = \sum_{j=1}^L N_j, \quad (2.35)$$

- where
- i is the nuclide number in the spectrum of energy-releasing nuclei;
 - j is the nuclide number in the spectrum of energy-consuming nuclei being products of nuclear combustion;
 - K is the total number of types of nuclides in the spectrum of nuclei that release energy;

- L is the total number of types of nuclides in the spectrum of nuclei that consume energy;
- $\Delta\bar{E}_i$ is the weighted average increase in the specific binding energy of nuclei that release energy as compared to the source nuclei of the collapsed matter;
- $\Delta\bar{E}_j$ is the weighted average decrease in the specific binding energy of nuclei that consume energy as compared to the source nuclei of the collapsed matter.

In the other extreme case of the scenario of the collapse self-development, where the maximum values of A for the created nuclides reach the area of superheavy nuclei, while the specific binding energy per nucleon approaches the Migdal limit (see Ref. 22) ($\bar{E} \approx 20$ to 40 MeV/nucleon), the following relation is valid:

$$\sum_{i=1}^K \Delta\bar{E}_i \cdot N_i \gg \sum_{j=1}^L \Delta\bar{E}_j \cdot N_j, \quad (2.36)$$

i.e., a significant free energy is released in the form of both the kinetic energy of transmuted nuclei and other particles produced by the nuclear combustion and as quanta of electromagnetic emissions in a wide frequency range.

In the last case (i.e., when Eq. 2.36 is true), the next stage of the process will develop.

Stage 4. Explosion of the cluster fragmentation products

- Explosive release of high-energy products of the collapse from HD, i.e., from the area where the collapsing wave-shell decelerates, stops, and explodes.
- Interaction of cluster explosion products with the nuclei evaporated from the wave trailing edge at the implosive stage of the process, as well as with the nontransmuted (i.e., source) matter that makes up the outer part of the volume of the target (i.e., of the collapsed body), where the wave-shell emerged and evolved before the first stage of the nuclear combustion began.

The nature and direction, in energy terms, of nuclear reactions that occur at the fourth stage are determined by the whole set of characteristic parameters of both the products of the explosion (fragmentation) of the created fireball, which are leaving HD area, as well as the transmuted and nontransmuted matter that surrounds this area at the moment of the collapse inversion. Consequently, the explosive stage of the process creates its

own energy balance, which, in its turn, may take the form of either the additional energy release at this explosive stage or the partial absorption of the free energy released at the first, i.e., implosive, stage of the process.

2.4. On the Technical Implementation, Choice of Driver Construction for Shock Compression, and Experimental Testing of the Effectiveness of Approach

Taking into account all required properties, as well as cost considerations, controllability, and potential ability to reproduce the key physical parameters of the desired driver, there has been no other choice but to use a high-current subrelativistic or relativistic electron beam.

In 1999, a small experimental facility was designed and put into operation at the Electrodynamics Laboratory “Proton-21”. It was a pulse vacuum-discharge facility with a total energy reserve of ca. 3 kJ and the efficiency of conversion of the accumulated energy into the electron beam energy up to 10%. The facility was aimed at the quick testing and the estimating of the efficiency of hypothetical mechanisms of shock compression of matter in solid-state targets, with the generation and, hopefully, the inertial and magnetic confinement (in the beam focal area) of a plasma bunch from the target material, with parameters that would allow one to satisfy and exceed the Lawson criterion (see Ref. 23) for inertial nuclear fusion.

It was expected that the beam driver will provide the required parameters for the formation of a superdense plasma bunch around the target focal point, with the following characteristics: $n \cdot \tau \approx 10^{15} \text{ s/cm}^{-3}$ (where n is density of the solid target, $n \sim 5 \cdot 10^{22} \text{ cm}^{-3}$, τ is duration of the pulse current (confinement time), $\tau \sim 10^{-7} \text{ s}$) and the ion temperature (for Cu targets) in the range of 3 to 10 keV.

In accordance with the criterion Eq. 2.30, in order to achieve the critical efficiency threshold, one has to minimize, with a limited energy reserve in the driver, the volume of the near-surface layer of the target which receives the most part of the shock impact energy and transfers it to the next thin layer of the target volume, thus creating a self-shrinking (self-collapsing), self-compressing wave-shell of the extreme density of energy and matter, scanning the target volume from its surface to the center (Fig. 2.3).

Having analyzed a number of known approaches, we have developed a method of providing the impact on a target (see Refs. 24, 25) through *bombarding its surface with a self-focusing electron beam* in the mode where it experiences the collective deceleration in the near-surface dense plasma and in the very thin surface layer of the target anode serving as energy concentrator, in order to excite a converging (collapsing) quasi-isoentropic solitary spherical wave of density of the highly ionized matter which then

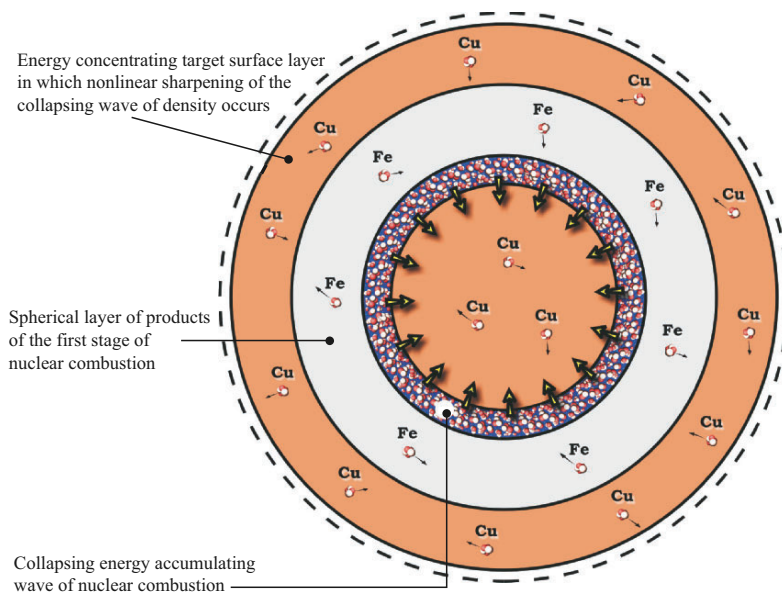


Fig. 2.3. Schematic picture of the intermediary stage of the nuclear combustion in a uniform copper (Cu) target.

shrinks to the target center, with the extreme state of the matter in the collapse microvolume.

The hard-current ($\sim 3 \cdot 10^4$ A) self-focusing electron beam was formed in a relativistic vacuum diode consisting of the metal anode serving as an energy concentrator, shaped as a cylinder conjugated with a hemisphere at the cathode-oriented edge, and the highly efficient explosive-emission plasma cathode providing the delivery of the beam of the pulse power of $\sim 10^{10}$ W to the anode surface with the power density of $\sim 10^{13}$ W $\cdot$ cm $^{-2}$ and the average pulse duration up to 10^{-7} s (see Refs. 24, 25).

The very first experiment using this arrangement, carried out on February 24, 2000, proved to be successful, by resulting in the explosion of the cylindrical target from inside with the creation of a crater passing into an axial channel (Fig. 2.4).

The nature of the damage meant that the maximum energy density was achieved precisely at the focus on the axis of the cylindrical target, and has been an indirect indication that the planned process took place.

When looking under the microscope at the exploded target, on the chemically pure surface of the copper-made accumulating screen which surrounded the bed of the concentrator target made of the same copper, we found a macroscopic area (around 1 mm long) of a solidified silver-and-white “lava” which had flowed out of the exploded “volcano” with a tubular

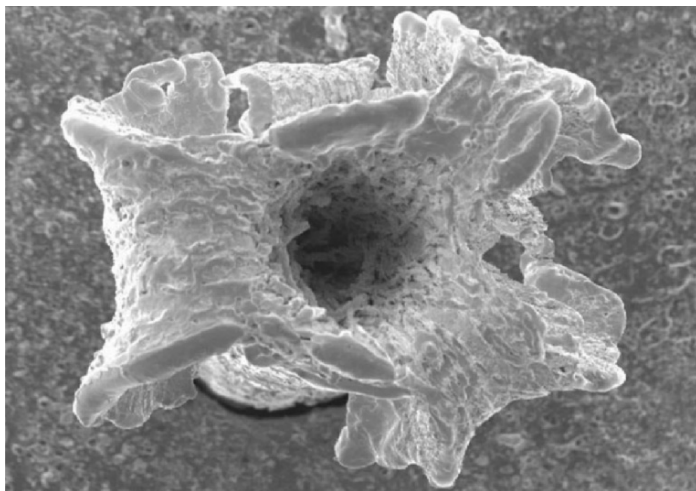


Fig. 2.4. Crater-like destruction of the initially uniform target anode that serves as the energy concentrator (from the face surface, cathode side) under the impact of the process initiated by an electron beam.

crater and left its traces, some drops, on the surface of one of the “petals” of the exploded tube which formerly was a uniform target rod.

The electron probe microanalysis of the element composition of that “lava” showed that it consists of zinc for 71%.

In 2000–2004, more than 20 000 analytical studies were carried out, using both physical and chemical methods, by many specialists of the Electrodynamics Laboratory “Proton-21”, as well as specialized analytical laboratories in Ukraine, Russia, USA, Germany, and Sweden. As a result, the following facts have been reliably established. The products released from the central area of the target destroyed by an extremely powerful explosion from inside in every case of the successful operation of the coherent beam driver created in the Electrodynamics Laboratory “Proton-21”, with the total energy reserve of 100 to 300 J, *contain significant quantities (the integral quantity being up to 10^{-4} g and more) of all known chemical elements, including the rarest ones. The local concentrations of those elements in various areas of the surface of accumulating screens made of chemically pure materials vary in the broadest range, from millesimals and centesimals of 50% to 70 % and more, with combinations, compounds, and alloys which cannot be obtained under usual conditions.* In addition, the ratios between the concentrations of stable isotopes of chemical elements in the thin near-surface layers of the matter produced as a result of the transformation of exploding targets, which are precipitated on accumulating screens, usually turn out to be significantly different from those found in the Nature. Explosion products

always contain such elements in significant amounts, including rare earths which were not found by any analytical methods, including ultra-sensitive mass-spectrometry, in the materials used in the details of the experimental setup, nor in those of the forming device.

Any attempts aimed at finding any *significant correlation between the concentrations* of chemical elements found on accumulating screens and the concentrations of the same elements in the impurities contained in materials used for the facility details contacting the vacuum chamber, where the exotic products are collected on the accumulating screen surface, were unsuccessful.

After five years of continuous research of various samples (accumulating screens), on which the matter was precipitated from exploding targets made of chemically pure metals (Mg, Al, Ti, V, Cr, Mn, Fe, Co, Ni, Cu, Zn, Nb, Mo, Pd, Ag, Ta, W, Pt, Au, Pb, Bi), it has been established that the correlation coefficients between the relative prevalence of chemical elements in the products of the target explosion and their concentration in impurities of the materials used in the construction of the setup and in the residual atmosphere of its vacuum system, vary from $+0.09$ to -0.17 , with the statistical average of -0.04 (with the confidence interval of 0.75). In other words, *there is no relation between the element composition of products of the target matter transformation and that of the initial target or materials used in manufacturing the setup details.*

Based on the data partially quoted above, as well as on the much larger amount of data quoted and justified in other parts of this book, the researchers, who have directly participated in the experiments and have become the authors of this book, consider it as the undoubtedly established experimental fact, that the products of the explosive transmutation of the target matter in experiments carried out at the Electrodynamics Laboratory “Proton-21” contain macroscopic quantities of artificially synthesized chemical elements or, in other words, products of artificial, or laboratory, nucleosynthesis.

The above facts had provided, in 2001–2003, the first grounds to say that, in the experiments carried out at the Electrodynamics Laboratory “Proton-21”, we have succeeded in implementing, and learning how to reliably reproduce under laboratory conditions, a microscopic analogue of the natural physical phenomenon which accounts for the explosive nucleosynthesis resulting in the creation of the full range of stable isotopes of all chemical elements found in the Nature and seems to be an energy source for supernova flares, as well as possibly a number of other astrophysical processes of the pulse nature.

This process, which can be reliably reproduced, is self-sufficient in energy terms. *The energy required to trigger the process is about 10^4 times*

lower than the whole work of the process on transforming the matter plus the total energy of the particle and radiation streams it produces.

The resulting products of the process (i.e., isotopes of the created chemical elements) are stable, irrespective of the activity of the target matter. This means that the discovered physical process, which can be controllably reproduced, can be used to create the efficient technologies of neutralizing the radioactive waste which would not consume much energy from external sources.

Having a basic physical process or any combination of such processes is not sufficient to create a real, safe, and profitable industrial technology. We have to learn how to maintain that process in the optimal mode, while keeping the full control over it. It is impossible to achieve this result without an adequate physical theory.

Obviously, in the situation in question, it was impossible to use an existing theory, due to the lack of such theories. A new one was needed.

Its creation, in its turn, needs a conceptual model or hypothesis, because [as Claude Bernard (see Ref. 26) put it] “A hypothesis is ... the obligatory starting point of all experimental reasoning. Without it no investigation would be possible, and one would learn nothing: one could only pile up barren observations. To experiment without a preconceived idea is to wander aimlessly.” In these terms, the experimental testing of the above-mentioned basic hypotheses on the mechanism of self-organizing and self-supporting nuclear combustion became a matter of principle.

Different parts of this book contain the sufficiently complete and detailed descriptions of not only the experimental results but also the methods of preparation and execution of experiments using the small experimental setup, IVR-1, carried out at the Electrodynamics Laboratory “Proton-21” within the framework of the Luch project in 2000–2004.

In this very special situation, the methodological information seems to be required, since it would be otherwise impossible to estimate the reliability and correctness of nontrivial statements made in this book. It enables interested experts to make sound proposals for the alternative methods of investigation and measurements, which would help to really understand the mechanism of this wonderful physical phenomenon, which has been discovered and is reproduced in a controlled way in the Kiev experiments on nucleosynthesis.

Electrodynamics Laboratory “Proton-21” would appreciate and consider, in a constructive way, any offers of cooperation in conducting the respective experiments using our experimental facilities.

3

EXPERIMENTAL SETUP

E. V. Bulyak, V. G. Artyukh, and A. S. Adamenko

The experiments aimed at the compression of a substance to superhigh densities are performed on a setup containing a generator of high-voltage pulses and a vacuum diode with needle-like electrodes.

The generator of high-voltage pulses exciting a beam in the vacuum diode must produce a pulse with a steep leading edge to transfer the beam rapidly and without losses to the quasipinch state.

It is desirable that a pulse have a descending trailing edge, which compensates the increase in the diode perveance arising at the expense of the dispersion of a cathodic plasma.

It turned out that a generator of high-voltage pulses with inductive accumulation of energy is most suitable for the spatio-temporal concentration of energy. Indeed, a generated pulse has a steep leading edge, whose duration is approximately equal to the switching time, and an exponentially descending trailing edge. The pulse of such a generator can feed the diode without intermediate forming.

The simplified electric circuit of the generator is given in Fig. 3.1; (see Refs. 28, 29). The generator operates in the following way. The preliminarily charged capacitor C is connected through an additional commutator (not shown) to the inductive storage unit L sequentially connected with the opening switch R . If the discharging current attains the first maximum, the fast opening switch (OS) increases its own resistance many times. In this case, the discharging current of the contour, which is maintained by the inductance, creates a voltage pulse on the opening switch used for feeding the load.

3.1. Generator Performance

The balance equation for a charge in the circuit of the system drawn in Fig. 3.1 reads

$$L \frac{\Delta^2 q}{\Delta t^2} + R \frac{\Delta q}{\Delta t} + \frac{q}{C} = 0, \quad (3.1)$$

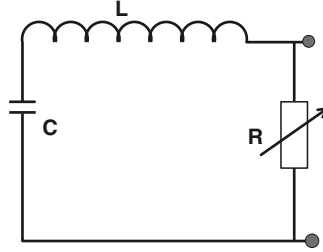


Fig. 3.1. Simplified scheme of the generator of high-voltage pulses with inductive storage unit.

where L is the inductance of the system (the inductive storage unit), C is the capacitance of the system (a storage capacitor), and R is the ohmic resistance of the system (the summary resistance of OS and the load connected in parallel).

The current in the system is given by

$$I = \frac{\Delta q}{\Delta t}$$

and the output voltage by

$$U = R \frac{\Delta q}{\Delta t} = IR.$$

In the time interval from the closing of the commutator connecting the charged capacitor with the discharging circuit to the moment when OS comes into action, we consider the resistance to be constant, $R = R_0$. Then the current in the system is given by

$$I(t) = -q_0 \exp(-\beta t)(\beta \cos \omega t + \omega \sin \omega t), \quad (3.2)$$

where

$$\omega^2 = \frac{1}{CL} \left(1 - \frac{R^2}{4L} \right) > 0, \quad \beta := \frac{R}{2L}.$$

We assume that OS comes into action in a quarter of the period, $\omega t_* = \pi/2$, when the capacitor is completely discharged. Then, at the actuation time $t = t_*$, the current in the system has the value

$$I_* = -q_0 \omega \exp(-\beta \pi/2\omega) = -CU_0 \exp(-\beta \pi/2\omega). \quad (3.3)$$

and the output voltage at the same time is

$$U_*^{(\text{os})} = RI_*.$$

Let us consider the idealized case where OS comes into action for an arbitrary small time. At the time $t = t_*$ (in a quarter of the period), let the resistance of OS (and the load connected in parallel) change stepwise from R_0 to R_1 . That is, a temporal behavior of the resistance has the step form

$$R(t) = R_0 + (R_1 - R_0) H(t - t_*), \quad (3.4)$$

where $H(x)$ is the Heaviside function.

In a time interval short compared with the period of oscillations of the system and with zero charge on the capacitor, Eq. 3.1 is reduced to the following equation for the current:

$$\frac{\Delta I}{\Delta t} + \frac{R(t)}{L} I = 0. \quad (3.5)$$

With regard to Eq. (3.4), its solution is

$$I(t) = I_* \exp \left[-\frac{(t - t_*)}{L} [R_0 + (R_1 - R_0) H(t - t_*)] \right]. \quad (3.6)$$

It follows from Eq. 3.6 that the current in the circuit is not changed at once after switching ($t = t_* + 0$) by virtue of the circuit lag caused by the presence of the inductance; the voltage across OS increases in proportion to the resistance ratio R_1/R_0 :

$$U_{*(+)} = U_{*(-)} R_1/R_0 = R_1 C U_0 \omega \exp(-\pi\beta/2\omega). \quad (3.7)$$

This formula expresses the maximum value of the output voltage in terms of the system parameters. Then Eq. 3.7 takes the form

$$U_* = U_0 R_1 \sqrt{\frac{C}{L} \left(1 - \frac{C R_0^2}{4L} \right)} \exp \left[-\frac{\pi R_0 C}{4\sqrt{CL - (C R_0/2)^2}} \right]. \quad (3.8)$$

As seen, at the infinitely large rate of switching, the maximum voltage is proportional to the product of the charging voltage by a finite resistance of OS. The dependence on the remaining parameters of the system (the initial resistance of OS, capacitance, and inductance) is nonlinear. In particular, Eq. 3.8 shows that, at zero initial resistance of the system, $R_0 \rightarrow 0$, the maximum voltage

$$U_{\max} = U_0 R_1 \sqrt{C/L} \quad (3.9)$$

is proportional to the root of the ratio of the storage capacitor to the system inductance.

If we define the system reactance R_r as

$$R_r := \sqrt{L/C},$$

then Eqs. 3.8 and 3.9 take a “dimensionless” form

$$U_* = U_0 \frac{R_1}{R_r} \sqrt{1 - \kappa_0^2} \exp \left[-\pi R_0 / 4 R_r \sqrt{1 - \kappa_0^2} \right], \quad (3.10)$$

$$U_{\max} = U_0 \frac{R_1}{R_r}, \quad \text{at } R_0 = 0, \quad (3.11)$$

$$\kappa_0 := \frac{R_0}{2R_r}.$$

Taking the current according to Eq. 3.6 and the zero charge in the capacitor as new initial conditions and assuming that the system resistance is equal to R_1 (the finite summary resistance of “OS + load”), we will find the solution of Eq. 3.1. By using the sewing condition for the current at the breaking time, we get the dependence of the output voltage on time at $\kappa_1 < 1$:

$$U_1(t) = U_{*+} \frac{\sqrt{1 - \kappa_1^2} \cos \omega_1 t - \kappa_1 \sin \omega_1 t}{\sqrt{1 - \kappa_1^2}} \exp(-\kappa_1 \omega_0 t). \quad (3.12)$$

In this case, the character of the dependence of the voltage on time is not changed; as before, the voltage is sinusoidal with the exponentially decreasing amplitude. After breaking, the frequency of oscillations decreases, and the decrement increases.

At the boundary of the oscillatory mode, $\kappa_1 = 1$, we have the exponential decay

$$U_1(t) = U_{*+} \exp(-\omega_0 t) = U_{*+} \exp \left(-t / \sqrt{CL} \right). \quad (3.13)$$

Upon the increase of the finite resistance of OS beyond the boundary of a quasiperiodic solution ($\kappa_1 = 1$), the form of a pulse becomes

$$U_+(t) = U_* \frac{\sqrt{\kappa_1^2 - 1} \cosh(\omega_1 \kappa_1 t) - \kappa_1 \sinh(\omega_1 \kappa_1 t)}{\sqrt{\kappa_1^2 - 1}} \exp \left(-t \omega_0 \sqrt{\kappa_1^2 - 1} \right). \quad (3.14)$$

It is seen from this formula that, shortly after the switching ($\omega_0 t \lesssim 1$), the voltage varies as

$$U_1(t \lesssim 1/\omega_0) \approx U_*(1 - 2\kappa_1 \omega_0 t), \quad (3.15)$$

i.e., the pulse is of a triangular form.

As seen from Eq. 3.15, *the voltage decay rate is proportional to a finite resistance of OS.*

By setting a certain level of the “valid” pulse, we can estimate its duration. For example, the pulse decays to the level η for the time

$$\tau = \frac{1 - \eta}{2\kappa_1\omega_0} = (1 - \eta) \frac{L}{R_1}. \quad (3.16)$$

The analytic consideration of the idealized scheme allows us to define the main parameters of the output high-voltage pulse.

We use the following notations:

L — inductance of the system, H;

C — capacitance of the storage capacitor, F;

U_0 — voltage across the charged capacitor, V;

R_0, R_1 — summary resistance of a load and OS connected in parallel in the closed and open states, respectively.

The voltage multiplication coefficient λ (the ratio of the maximum output voltage to the voltage across the charged capacitor) is defined as

$$\lambda = \frac{U_{\max}}{U_0} = \frac{R_1}{R_r} \sqrt{1 - \kappa_0^2} \exp\left(-\frac{\pi\kappa_0}{2\sqrt{1 - \kappa_0^2}}\right), \quad (3.17)$$

$$R_r := \sqrt{L/C}, \quad \kappa_0 := R_0/2R_r.$$

The form of a pulse of the output voltage is defined by the formula

$$U_+(t) = \frac{U_{\max} \exp\left(-t\omega_0\sqrt{\kappa_1^2 - 1}\right)}{\sqrt{\kappa_1^2 - 1}} \left[\sqrt{\kappa_1^2 - 1} \cosh(t\omega_0\kappa_1) - \kappa_1 \sinh(t\omega_0\kappa_1) \right], \quad (3.18)$$

where $\omega_0 = 1/\sqrt{LC}$, and the time is counted from the actuation time of OS.

The width (duration) of a high-voltage pulse on a height η relative to the maximum value taken as 1 is equal to

$$\tau = (1 - \eta) L/R_1. \quad (3.19)$$

The main condition for the enhancement of the output voltage (as compared to the charging voltage of the capacitor) and a time shortening of a pulse is a jumplike increase in the summary resistance of the load and OS connected in parallel to a value exceeding the circuit reactance.

It is worth noting that, upon the actuation of OS at the current maximum when the voltage across the capacitor is zero, the output voltage is equal to the voltage across the inductance:

$$U = L \frac{\Delta I}{\Delta t}.$$

This can be used to estimate the output voltage by the decay rate of the current.

3.2. Numerical Model of the Setup

The analytic model of the generator considered in the previous section allowed us to derive the main characteristics of the output pulse. However, the real setup is a much more complicated system. Among such “complications”, most essential are the following ones:

- the load is a diode, whose impedance depends on the applied voltage;
- the diode is connected through the inductance, i.e., the load is not purely active;
- the plasma opening switch (POS) comes into action not obligatory at the peak value of the current and for a finite time;
- at the expense of the movement of the plasma strap (upon the use of POS), the storage inductance varies (increases) during the stage of storage.

This system does not yield to theoretical analysis. Therefore, we developed a numerical model, code SAEB–2002. Within the model schematically presented in Fig. 3.2, we numerically solved a more general equation, by representing the first term in Eq. 3.1 in the form that takes into account the dependence of the storage inductance on time:

$$L \frac{\Delta^2 q}{\Delta t^2} \rightarrow \frac{\Delta}{\Delta t} L \frac{\Delta q}{\Delta t}.$$

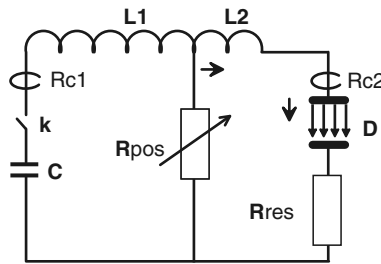


Fig. 3.2. Electrical scheme as a basis of the numerical model. k – key; D – diode; Rc1, Rc2 – integrating Rogowski coils; L1, L2 – storage and load inductances ($L1 + L2 = \text{const.}$); Rpos – resistance of POS; Rres – residual resistance of the diode. The bold arrows show the directions of motion of the plasma strap up to the breaking time and the cathodic plasma of the vacuum diode.

A breaking of POS occurs for a finite given *breaking time* τ_{sw} after the termination of a prescribed time interval t_{bsw} counted from the starting time of the system (the actuation of the commutator “k”):

$$R_{\text{pos}} = R_0 \quad t \leq t_{\text{bsw}}, \quad (3.20)$$

$$R_{\text{pos}} = R_1 \quad t \geq t_{\text{bsw}} + \tau_{\text{sw}}. \quad (3.21)$$

The (planar) vacuum diode serves as a load, whose ohmic resistance depends on the voltage applied to it according to the well-known “3/2” law:

$$R_d = 1/\kappa\sqrt{U};$$

here, κ is the diode perveance varying in time at the expense of the dispersion of the cathodic plasma in accordance with

$$\kappa(t) = \kappa_{\text{in}} \frac{d^2}{(d - v_{\text{pl}}t)^2};$$

where v_{pl} is the given speed of dispersion of the plasma and κ_{in} the initial perveance of the diode.

For a more adequate comparison of the numerical model to the real setup, we add the former by two integrating Rogowski coils which were modeled with equations given in Ref. 30:

$$\frac{\Delta I_{\text{Rc}}}{\Delta t} = -\frac{R_{\text{Rc}}}{L_{\text{Rc}}} I_{\text{Rc}} + \frac{\Delta I}{\Delta t}, \quad U_{\text{Rc}} = g R_{\text{Rc}} I_{\text{Rc}}. \quad (3.22)$$

Here, U_{Rc} is the output voltage of a Rogowski coil, and g the gauge coefficient.

The “oscillograms” derived within the numerical model turned out to be qualitatively and quantitatively similar to real ones.

The study of the numerical model showed that the maximum voltage of a high-voltage pulse is significantly defined by a value of the load at the open POS. This voltage depends quite weakly on the switching time (of course, in the case where the switching time is considerably less than the period of natural oscillations of the system with the closed POS).

An example of model oscillograms is presented in Fig. 3.3.

It is seen from this figure that a pulse consists of three sections: The first is periodic and corresponds to the storage of energy in the inductive storage unit; the second is aperiodic and associated with the actuation of OS and the generation of a high-voltage pulse; the third section is related to the dissipation of the energy remaining in the system and is omitted from consideration.

The second section, a high-voltage pulse proper, is shown in Fig. 3.4.

It is seen from this figure that model signals from integrating Rogowski coils repeat, on the whole, real current pulses.

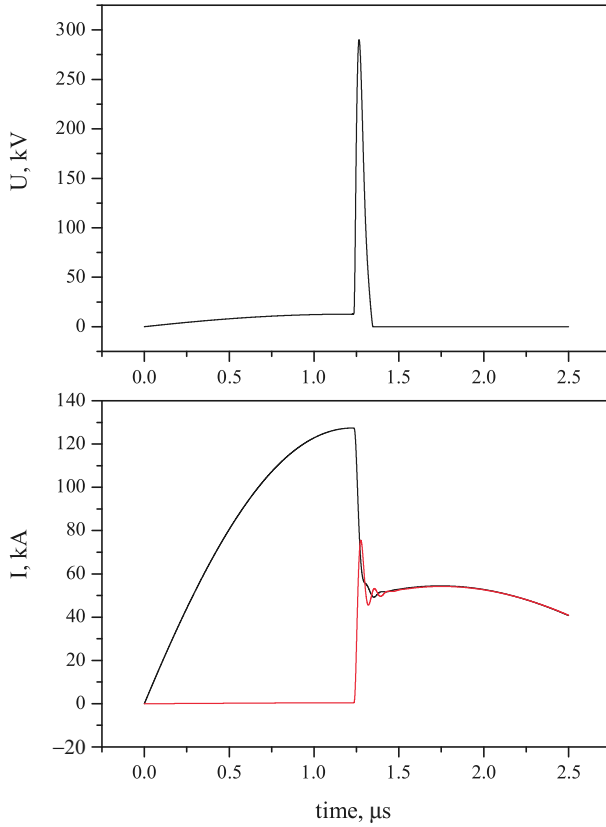


Fig. 3.3. Lower figure—current pulses in the generator and diode (the grey line). Upper figure—voltage pulse across the diode.

3.3. Construction of the “Pulse Generator—Vacuum Diode” System

In order to actuate an energy-concentrating relativistic vacuum diode (RVD), we manufactured a generator of high-voltage pulses with a high peak power which contains a plasma opening switch of current (POS); (see Refs. 30, 31).

The scheme of the system is given in Fig. 3.5. A spark gap connects a capacitor of $3\mu\text{F}$, 50 kV with the OS cathode mounted on a high-voltage insulator. In the experiments with currents of the electron beam of more than 70 kA, we used a Marx bank with a stored energy of 75 kJ instead of a capacitor and a spark gap. The POS and its load (RVD) were positioned in a grounded vacuum chamber (anode). In this case, the cathode and anode of POS form a vacuum coaxial. From 6 to 12 plasma sources are used for the injection of a plasma through a steel mesh towards to the

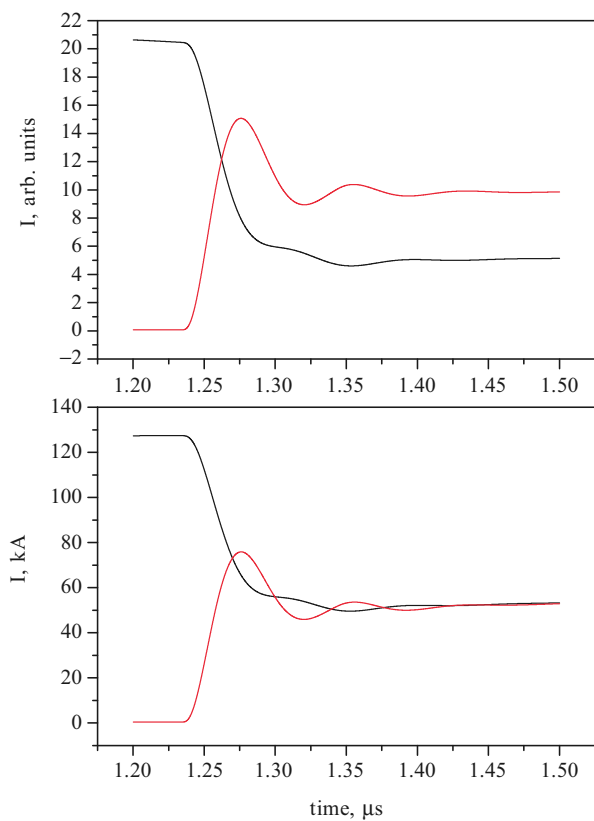


Fig. 3.4. Lower figure—current pulses in the generator and diode (the grey line). Upper figure—corresponding signals from Rogowski coils.

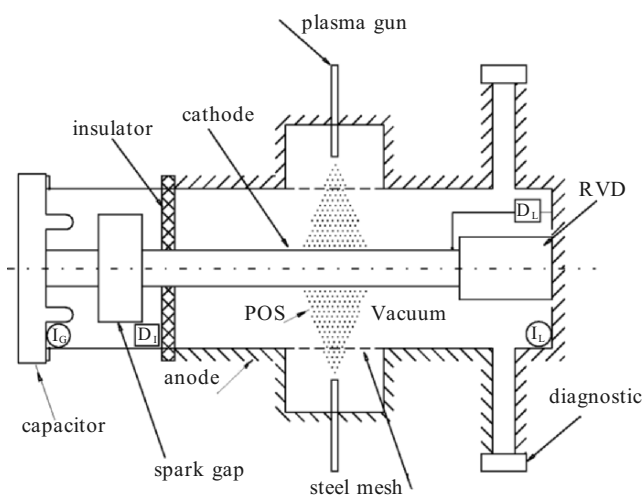


Fig. 3.5. Construction of the generator.

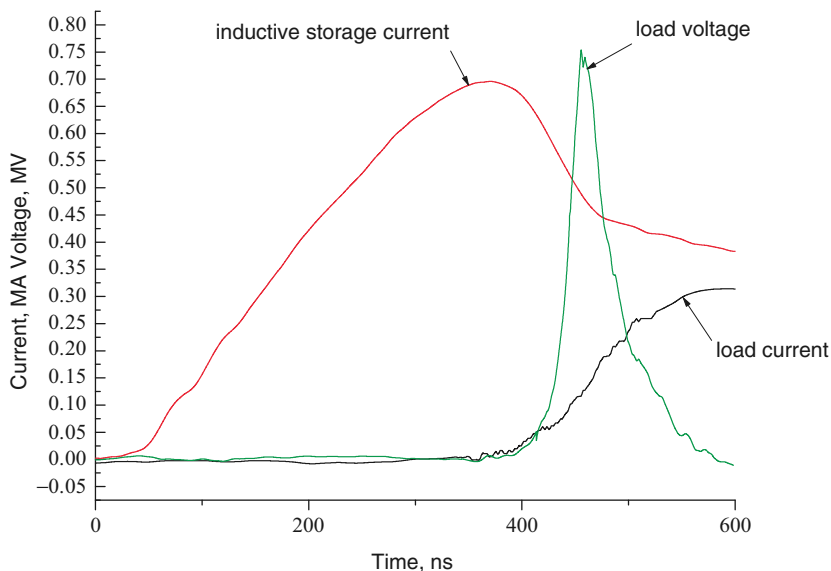


Fig. 3.6. The typical oscillograms of the currents of the inductive storage unit and the vacuum diode, as well as the voltage across the diode.

central conductor of the coaxial (cathode). The applied sources of a plasma are close by construction to cable guns described in Ref. 32.

The action of POS begins with the running of the discharging current of the capacitor through the plasma preliminarily injected by plasma guns into the region between the anode and cathode. When this current reaches a threshold value, there occurs a fast increase in the impedance of the volume of a plasma in POS. As a result, the energy stored by this time moment in the magnetic field of the inductance of the discharge contour (in the inductive storage unit) is spent for the maintenance of current in the contour, and the increased voltage across POS is applied to the diode through a transmission line (a section of the vacuum coaxial positioned between POS and RVD). The appearance of an electron beam in the diode leads to the branching of a part of the current of POS into the diode and to the release of a part of the stored energy there.

As usual, the resistance of the open POS is equal to several Ohm. At the current in the inductive storage unit of 150 to 1000 kA at the breaking moment, the voltage across RVD attains 0.3 to 1.5 MV.

The currents running in the system were measured with shielded Rogowski coils (see Ref. 30) positioned on the capacitor (I_G) and on the load (I_L).

The voltage across POS is determined by means of a resistive voltage divider mounted on the high-voltage insulator on the side of the spark gap

(D_I), whose indications were subjected to the inductive correction with the use of current signals. The direct measurements of the voltage across RVD were carried out by using a resistive divider connected in parallel (D_L).

Prior to the mounting in the system, Rogowski coils and voltage dividers were calibrated.

In Fig. 3.6, we presented the typical oscillograms of the currents of the inductive storage unit and the vacuum diode, as well as the voltage across the diode.

3.4. Results and Conclusions

Thus, the generator of high-voltage pulses manufactured in the Electrodynamics Laboratory on the basis of an inductive storage unit with a plasma opening switch satisfies the posed requirements. Generated pulses possess a steep leading edge (about 10 ns). The form of a generated pulse is close to a triangular one with rounded angles.

The developed numerical model adequately describes the operation of the generator loaded by a vacuum diode.

The model calculations confirmed by measurements showed that the setup efficiency is sufficiently high and reaches 20–30%.

Part II

Some Experimental Results

OPTICAL EMISSION OF A HOT DOT (HD)

V. F. Prokopenko, A. I. Gulyas, and I. V. Skikevich

In the experiments performed on the setup IVR-1, the processes of destruction of the anode and dispersion of a plasma bunch are associated with a bright light flash. In studying the optical emission appearing at this stage a significant interest was paid to the information about the kinetic energies of separate ions and the total energy yield with the corpuscular component of a plasma bunch, as well as, to that about the intensity and spectral composition of a light pulse.

To derive the data on the kinetic energies of plasma particles, the time-of-flight method is ordinarily used (e.g., Ref. 161). The rough determination of the integral energy of a plasma bunch can be realized as a result of calorimetric measurements (see Ref. 162). However, the calorimetric measurements of energy yield do not guarantee the absolute inelasticity of the collisions of plasma particles with the surface of a calorimeter and, respectively, the reliability of the determination of the reflected part of a plasma bunch. In addition, it is necessary to take into account and exclude the energy released in a calorimeter upon the absorption of various emissions.

The information of interest to us can be derived by using the spectroscopic method of determination of the energy distributions of atoms excited in various processes (see Ref. 163). In this case, the necessary data can be derived from the Doppler broadening of spectral lines. If the broadening due to the Doppler effect is dominant, the information about the kinetic energies of emitters can be directly derived from the contour of a spectral line. In the general case, it is necessary to separate the Doppler component of the full contour. Just this method was used to derive the data on the kinetic energies of separate ions and the total energy yield. Below, we present the results of the performed studies.

4.1. Measuring Facilities

In the measurements of the optical emission of a plasma bunch in the vicinity of HD, we used a spectrometer SL 40-2-3648 USB, being a multichannel system of registration of optical spectra on the basis of charge transfer devices

(CTD). The spectrometer includes two spectrographs for the ranges of near-UV and visible light (216 to 700 nm), and for the IR range (677 to 1119 nm) with the focal distances of the camera objective to be 40 mm. A dispersive element in each spectrograph was the grating array with the following parameters: 600 lines per mm, inverse linear dispersion, respectively, 35.09 nm/mm and 31.95 nm/mm, resolution <1.5 nm in wavelength, and $12 \times 10 \text{ mm}^2$ in size.

The measurements were realized in the azimuthal plane of the setup through the output window made of Plexiglas of 4 mm in thickness. The transport of the optical emission from the vacuum chamber of the setup to a receiver occurs along a dielectric pipe of 75 mm in diameter. The receiver of light emission was located in measurements at the distance of ~ 6 m from HD and was connected by a doubled optical fiber cable of 2 m in length with the spectrometer and a high-speed photoelectronic multiplier FEU-30. Signals from FEU-30 were registered with a digital oscillograph with a transmission band of 1 GHz. Thus, in parallel with the measurement of the optical emission spectrum in the wavelength range 216 to 720 nm, we have registered the time signals from light flashes.

The calibration of a spectrometer was carried out:

- In the near-UV—by the continuous spectrum of a deuterium lamp DDS-80 with the certificate value of light intensity of $\sim 1.8 \text{ mW/sr}$ in the spectral range 215 to 300 nm;
- In the range of visible light—by the continuous spectrum of a tungsten filament lamp of 100 W in power.

The quality of the calibration in the range of optical emission was verified on separate lines of copper with known intensity.

4.2. Results of Measurements and Discussions

Prior to the presentation of the results of spectral measurements of the optical emission of HD derived with the use of the equipment described above, we give certain data concerning the photoregistration of the processes running in the vicinity of HD.

In Fig. 4.1a, we show the photo of a copper target of $500 \mu\text{m}$ in diameter before the experiment, and Fig. 4.1b presents the X-ray image of this target during the experiment which was obtained on an X-ray “KODAK” film with the help of an obscure-chamber. It is seen that no significant changes of the target sizes occur during the emission time of HD in the X-ray range (for 10 to 15 ns by the data of the measurement of X-ray pulses by GaAs detectors).

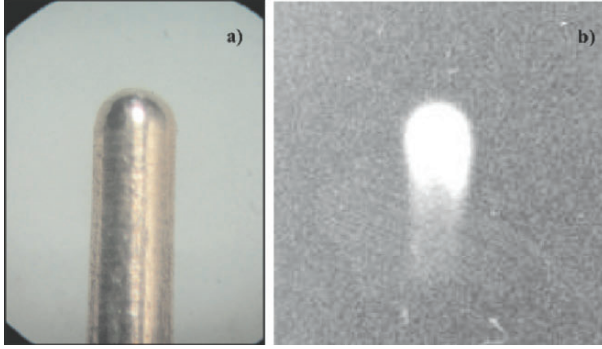


Fig. 4.1. Target before the experiment (a). The target material is copper (99.99 mass. % Cu). X-ray image of the target obtained with the use of an obscure-chamber on a “KODAK” film (b).

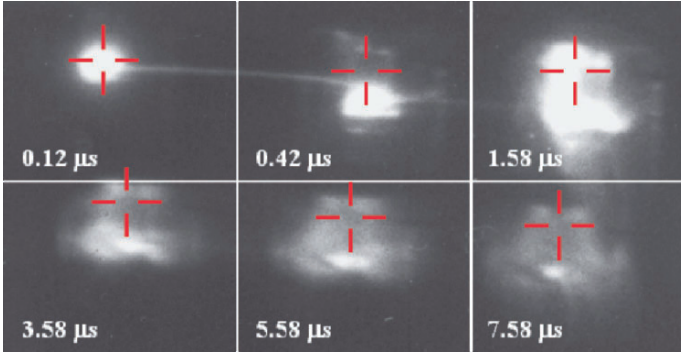


Fig. 4.2. Subsequent stages of the emission of a plasma bunch in the vicinity of HD.

Figure 4.2 illustrates the temporal stages of the emission of a plasma bunch in the optical range. The photos of the plasma bunch were obtained at various time moments from the beginning of the process (the commutation time of a current into the diode) with the help of a chronograph FER-7 modernized for the use in the “temporal lens” mode. The duration of separate frames was 50 ns. It is seen from the presented photos that the main information on the dispersion of a plasma bunch is contained, in fact, in the first frame with characteristic times of ~ 100 to 200 ns which are intrinsic to the process under study. This is testified also by the maximum amplitudes of signals which are registered by a fast photoelectric multiplier in the same temporal interval and thus characterize the largest brightness of a light flash. The remaining frames describe, apparently, the interaction of the dispersing plasma bunch with the substrate, which leads to the reflection of plasma particles and to the evaporation and dispersion of a substrate material.

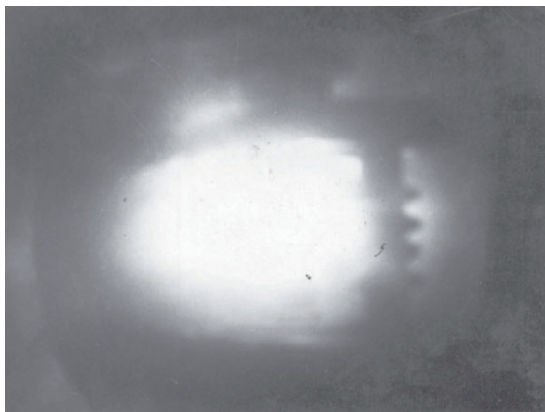


Fig. 4.3. Photo of the emission of a plasma bunch.

In Fig. 4.3, we present the photo of the luminescence of a plasma bunch, which was derived on a black-white photographic “Foto-64” film positioned at a distance of 23 cm from HD through a neutral light filter NS-10 and a lead glass with a photcamera “Smena”. Upon the survey, the lens diaphragm size was 1 mm. The exposure duration was practically equal to the duration of the luminescence of a plasma bunch.

The spectral measurements in the optical range were performed in the visible region and near-UV one in the integral mode without temporal resolution by individual lines. For such a nonstationary source as the studied plasma bunch, the task to attain both the temporal and spectral resolutions would require, apparently, an unconventional compromise upon the use of a spectral device and would generate significant technical difficulties. This follows also from the form of the optical spectra registered in experiments. An optical spectrum characteristic of the experiments with copper targets is given in Fig. 4.4.

As seen, the spectrum includes more than 300 emission lines unresolved by contour. The light energy yield in the optical spectrum derived by integration over all spectral lines in the wavelength range 300 to 700 nm can be equal to about several Joules. The brightness of an optical flash in the mentioned range of wavelengths is $\sim 5 \times 10^9 \text{ W}/(\text{m}^2 \cdot \text{sr})$ for the characteristic size of a plasma bunch of $\sim 3 \text{ cm}$ and the luminescence time of $\sim 100 \text{ ns}$.

In the derivation of the data on the energy of a plasma bunch, the essential point was the separation of individual spectral lines in the experimental spectra, since the contour of each spectral line is a source of the required information within the used method. The plasma parameters essentially influence the form and intensity of spectral lines and, most strongly, the line width.

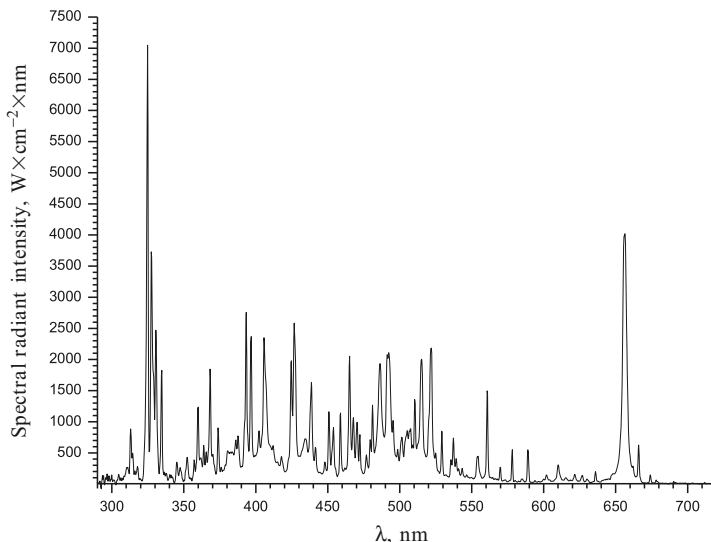


Fig. 4.4. Spectrum of the optical emission of a plasma bunch in the vicinity of HD (target: Cu).

The full width of a spectral line on the half-maximum intensity level appears as a result of the action of various broadening factors: the Doppler effect; Stark effect, the collective action of the electric microfields of charged particles on emitters; the emission attenuation effect as a result of the collisions of particles; the instrumental effects (the apparatus broadening), etc.

In plasma with the Doppler broadening of spectral lines, the Stark broadening is dominant. In addition, the additional broadening of lines (analogously to the Doppler broadening) can be induced by nonthermal motions: a macroscopic motion and microturbulence usually occurring in a nonstationary plasma of pulse discharges.

As usual, in the determination of the temperature of ions, it is essential to separate the purely temperature-related contribution to the broadening. In our case, this point is inessential, because the kinetic energy of dispersed ions of a plasma bunch was measured. Let us suppose that the observed profile of a spectral line is well described by the Gauss formula. Then we can surely consider the contributions of other mechanisms of broadening to be negligibly small, because they are associated with Lorentz profiles. In this connection, it is convenient to decompose every factor broadening a spectral line into two components: the Gauss term related to the Doppler broadening and the Lorentz term characterizing the emission attenuation as a result of collisions, the Stark effect, etc.

By assuming the statistical independence of the mechanisms of broadening of separate lines and by performing the convolution of the Gauss and Lorentz contours of these two components, we arrive at the well-known Voigt contour (see Ref. 164):

$$y(\lambda) = \frac{S_0 \cdot d}{\pi \cdot \sqrt{\pi} \cdot \sigma^2} \cdot \int_{-\infty}^{\infty} \frac{e^{-t^2} dt}{\frac{d^2}{\sigma^2} + \left(\frac{\lambda - \lambda_0}{\sigma} - t \right)^2}, \quad (4.1)$$

where λ is wavelength, S_0 is area of a spectral line, λ_0 is wavelength of a spectral line, d is half-width at the half-height of a Lorentz contour, σ is half-width on the $1/e$ -level of a Gauss contour.

The processing of spectra of the optical emission of HD was carried out with the use of the standard software “Peak Fit” allowing one to fit line spectra including up to one hundred of spectral lines. Because the registered optical spectra included a significantly greater number of lines, the analyzed spectrum was divided into several fragments for their processing. The processing of each fragment of a spectrum was performed separately.

While processing the measured optical spectra, we used a Voigt contour in the above-presented form. First, we eliminated the contribution introduced by the apparatus function to the profiles of spectral lines. The data on the apparatus function of the device were derived as a result of measurements of the spectrum of a mercury lamp. Then we decomposed the measured spectrum into components, i.e., into separate spectral lines. During the processing, we varied the width and form of a line (the ratio d/σ) up to the derivation of the optimum agreement between the total contour of spectral lines and experimental points. During the processing of spectra, we used a number of statistical criteria set in the processing program to control the spectrum decomposition exactness.

In Fig. 4.5, we show one of the fragments of the processed spectrum with the purpose to demonstrate the fitting quality. The lines in spectra were identified by using the tables of spectral lines (see Ref. 165).

We estimated the energy of the ion component of a plasma bunch as follows. In their essence, the Gauss components of spectral lines found due to the decomposition of the spectrum are the distributions over the velocity projections of emitting atoms on the observation direction, because

$$\Delta\lambda = \lambda v/c, \quad (4.2)$$

where λ is wavelength, v is emitter velocity, c is light velocity.

The distribution over the velocities of ions for isotropic angular distributions of dispersed particles can be obtained from the relation

$$\phi(V_x) = \int_{V_x}^{\infty} (1/v) P(v) dv, \quad (4.3)$$

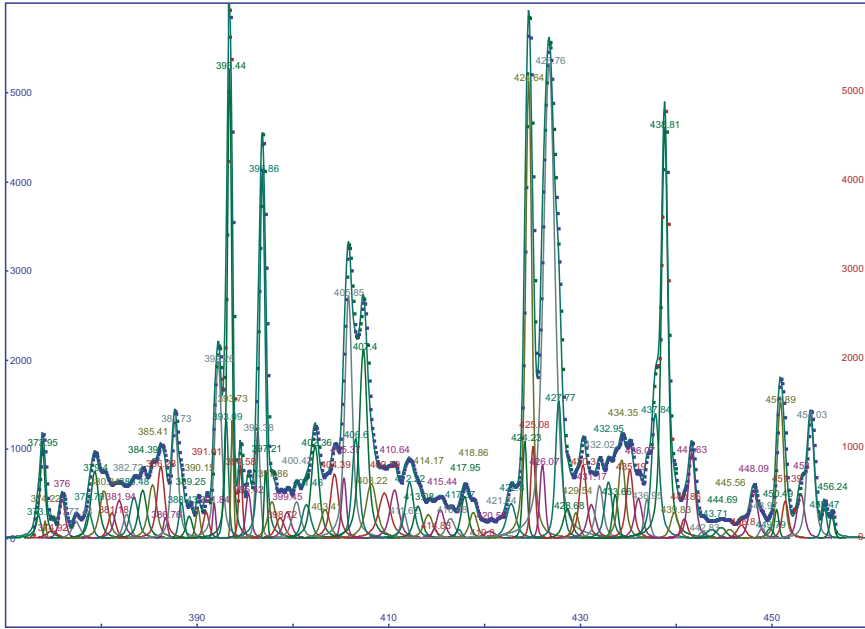


Fig. 4.5. Fragment of the spectrum of the optical emission of a plasma bunch in the vicinity of HD processed with the help of the program “Peak Fit”.

where $\phi(V_x)$ is distribution of the intensity over the contour of a spectral line; $P(v)$ is distribution over the velocities of emitting ions.

The above-presented relation yields the distribution over velocities of the emitting ions, $P(v)$, by differentiation.

To derive the distribution over velocities for the ions of separate chemical elements, we took the sum of Gauss components of all the spectral lines of a certain element with the parameters obtained as a result of the mathematical processing of the spectrum and with the weight coefficients expressed through the relative amount of ions emitting on each separate line on the left-hand side of the integral equation.

To estimate the number of emitting ions, we used the plots of the integral absorption of spectral lines (see Ref. 166). The distributions over the energies of ions of a plasma bunch were obtained from the corresponding distributions over velocities by a change of variables. Further, the distributions over velocities and energies were used to derive the estimates of the mean velocities and energies of ions of the plasma bunch.

As an example, Figs. 4.6 and 4.7 give the distributions over velocities and energies of hydrogen ions in the experiment with a copper target. In calculations, the intensity of separate spectral lines was normed on the effective cross-section of a plasma bunch and the emission duration. The

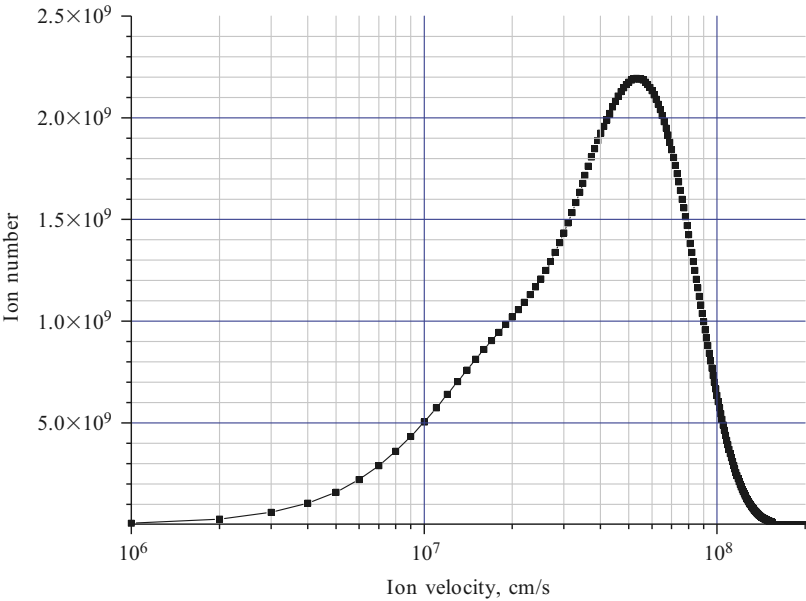


Fig. 4.6. Distribution over the velocities of H ions of a plasma bunch in the vicinity of HD (exper. No. 6466, 09 Feb 04).

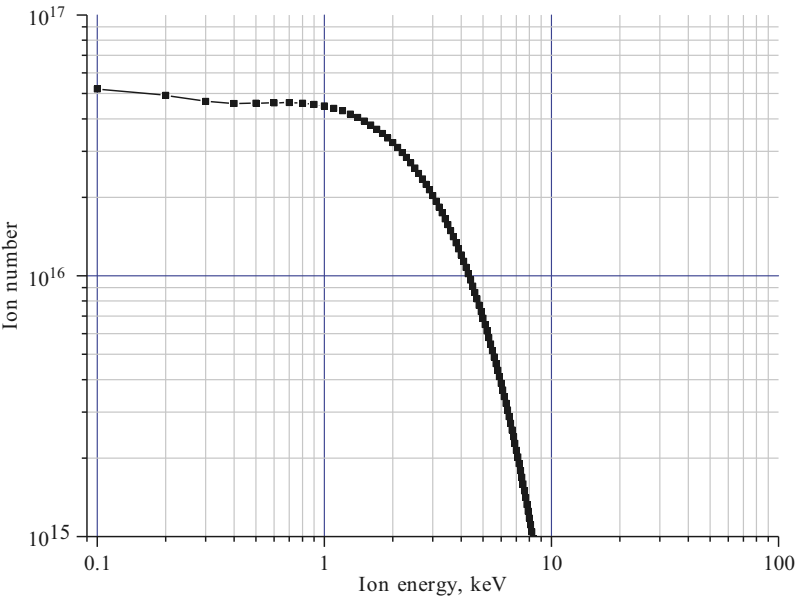


Fig. 4.7. Distribution over the energies of H ions of a plasma bunch in the vicinity of HD (exper. No. 6466, 09 Feb 04).

last was taken equal to the duration to a light flash signal on the half-height and was, as noted above, of the order of 80 to 100 ns. The effective sizes of a plasma bunch, being of the order of 3 cm by the data of the photoregistration, are defined by the mean velocities of ions and the time interval to attain the plasma luminescence maximum (~ 50 ns in all experiments).

To estimate the electron density of a plasma bunch, we used the Lorentz components of spectral lines by assuming that the broadening of spectral lines occurs due to the Stark effect under the action of electrons. While deriving the estimations, we used the formula of impact theory (see Ref. 164) applied in wide intervals of electron densities and temperatures:

$$\Delta\lambda_{1/2} \approx 2 \left[1 + 1.75 \cdot 10^{-4} N_e^{1/4} \alpha \left(1 - 0.068 N_e^{1/6} T^{-1/2} \right) \right] 10^{-16} \omega N_e, \quad (4.4)$$

where $\Delta\lambda_{1/2}$ is full half-width of the Lorentz component of an atomic line; N_e is electron density; T_e is temperature of electrons in $^\circ\text{K}$; α is parameter characterizing the broadening of a line due to the interaction with ions; 2ω is half-width of a line conditioned by electron impacts (at $N_e = 10^{16} \text{ cm}^{-3}$).

This approximate formula is sufficiently exact if the following inequality is valid:

$$10^{-4} N_e^{1/4} \alpha < 0.5. \quad (4.5)$$

For the application of the formula to singly ionized atoms, it is necessary to substitute the numerical factor 0.068 by 0.11. The tables of the parameters of α and 2ω for many main lines of neutral and single-ionized atoms of various elements from He to Ca and of Cs are given in Ref. 167. From the position of practical applications, the data of the tables cover the temperature interval from 5000 to 40 000 K. We will illustrate their importance by the example of hydrogen. At temperatures lower than the above-presented, hydrogen exists mainly in the form of molecules and is fully ionized at higher temperatures. In these limiting cases, the intensity of the atomic lines of hydrogen is too low to be registered. In our experiments, the line Ha is one of the most intense ones. Therefore, in view of the integral character of our measurements (without separation of the spectral lines of individual elements during the luminescence), we may consider the electron temperature not exceeding, in any case, 100 eV at the stage of the plasma bunch luminescence.

The other source for the derivation of estimates of the electron density is the relation of Inglis–Teller (see Ref. 166) which is especially convenient in the study of fast processes:

$$\lg N_e = 23.46 - 7.5 \lg n_{\max}, \quad (4.6)$$

where N_e is electron density; $n_{\max}$ is quantum number of a last visible line of the Balmer series of hydrogen.

Table 4.1. Estimates of the plasma bunch power in the experiments with targets and substrates made of different materials.

Expe- riment	Tar- get	Subs- trate	Electron density, cm^{-3}	Plasma bunch V, cm^3	Number of elec- trons	E_e , J	Number of emit- ting ions	E_i , J
6427	Pb	Cu	1.49E+17	9.6	1.43E+18	23	6.14E+17	2.08E+03
8140	Pb	Ta	1.76E+17	6.2	1.09E+18	17	2.43E+17	6.87E+02
8190	Pb	Al	8.26E+16	10.5	8.67E+17	14	1.80E+17	4.82E+02
8197	Pb	—	1.13E+17	12.6	1.42E+18	23	1.80E+17	5.75E+02
6466	Cu	Cu	1.86E+17	12	2.23E+18	38	6.14E+17	1.42E+03
6759	Cu	Cu	1.55E+17	9.8	1.52E+18	24	4.17E+17	1.47E+03
8171	Cu	Ta	1.62E+17	6.5	1.05E+18	17	2.80E+17	5.37E+02
6100	Al	Cu	1.50E+17	11.2	1.68E+18	27	4.69E+17	1.11E+03
8134	Al	Ta	7.46E+16	8.4	6.27E+17	10	1.77E+17	261

This relation is valid for hydrogen at temperatures above $T = 10^5/n_{\max}$ (K). At lower temperatures, a derived value of the electron density should be halved.

With the use of the described procedure of calculations, we derived the estimates of the plasma bunch power in experiments with different materials of the target and substrate. In Table 4.1, we give the estimates of the electron density, total numbers of electrons and ions in the plasma bunch volume, and energies of the electronic and ionic components. The electron density in each experiment was estimated by several spectral lines of elements C, O, He, Li, Na, and others, for which we succeeded to find the Stark broadening coefficients in the above-indicated sources. The results of estimates were averaged by several lines for each experiment.

The estimate of the electron density in the experiments by the Inglis-Teller relation (we identified five hydrogen lines of the Balmer series, $n_{\max} = 7$, in the optical spectra) gave $1.3 \times 10^{17} \text{ cm}^{-3}$, i.e., this value is also close to the derived estimates.

We note that the coincidence of the estimates derived by two different methods indicates, to a certain extent, the consistency of the results.

In Table 4.2, we present the detailed information on the element composition of a plasma bunch, mean velocities and mean energies of the ions of various chemical elements, number of emitting ions, and energy yield of the ionic component for various chemical elements in the experiment with a Pb target and a Cu substrate (exper. No. 6427). Table 4.2 contains also the composition of admixtures in the anode material.

Table 4.2. Estimates of the energy yield of the ionic component of a plasma bunch in the experiment with a Pb target and a Cu substrate (Experiment No. 6427).

	$\bar{E}$, keV	$\bar{v}$, cm/s	Emitting atoms		Energy yield		Ads*
			$N_{\text{em.at.}}$	%	E , J	%	
H	3.1E+00	7.7E+07	8.0E+16	1.3E+01	4.0E+01	1.9E+00	6.8E+14
He	3.1E+00	3.8E+07	8.4E+15	1.4E+00	4.1E+00	2.0E-01	
C	5.8E+00	3.1E+07	2.4E+16	3.9E+00	2.2E+01	1.1E+00	3.4E+14
N	5.3E+00	2.7E+07	2.6E+16	4.2E+00	2.2E+01	1.1E+00	5.6E+14
O	8.5E+00	3.2E+07	3.2E+16	5.2E+00	4.4E+01	2.1E+00	5.0E+14
F	6.6E+00	2.6E+07	2.9E+15	4.7E-01	3.0E+00	1.5E-01	7.7E+13
Ne	1.1E+01	3.2E+07	4.1E+15	6.7E-01	7.1E+00	3.4E-01	
Na	3.8E+00	1.8E+07	3.0E+15	4.9E-01	1.9E+00	8.9E-02	4.7E+16
Mg	1.7E+00	1.2E+07	4.0E+14	6.5E-02	1.1E-01	5.1E-03	9.3E+15
Al	7.8E+00	2.4E+07	1.2E+16	2.0E+00	1.5E+01	7.2E-01	5.7E+16
Si	2.5E+00	1.3E+07	1.4E+16	2.2E+00	5.4E+00	2.6E-01	5.9E+16
P	6.7E+00	2.0E+07	1.4E+16	2.3E+00	1.5E+01	7.1E-01	1.5E+15
S	5.9E+00	1.9E+07	1.6E+16	2.7E+00	1.5E+01	7.4E-01	2.2E+16
Cl	1.3E+01	2.7E+07	1.1E+16	1.9E+00	2.4E+01	1.1E+00	2.3E+16
K	1.2E+01	2.4E+07	2.2E+16	3.5E+00	4.2E+01	2.0E+00	3.4E+13
Ca	2.0E+01	3.1E+07	2.5E+16	4.0E+00	7.9E+01	3.8E+00	9.7E+15
Ti							1.9E+15
V	1.8E+01	2.6E+07	7.0E+15	1.1E+00	2.0E+01	9.7E-01	1.0E+14
Cr	9.9E+00	1.9E+07	1.0E+15	1.6E-01	1.6E+00	7.7E-02	1.8E+15
Mn	9.4E+00	1.8E+07	1.2E+16	2.0E+00	1.8E+01	8.8E-01	3.2E+14
Fe	1.7E+01	2.4E+07	6.8E+16	1.1E+01	1.9E+02	9.0E+00	5.1E+16
Co	2.6E+01	2.9E+07	8.8E+14	1.4E-01	3.6E+00	1.7E-01	4.3E+13
Ni	4.8E+01	4.0E+07	7.1E+14	1.2E-01	5.5E+00	2.6E-01	1.2E+15
Cu	3.2E+01	3.1E+07	8.6E+16	1.4E+01	4.4E+02	2.1E+01	3.3E+16
Zn	3.0E+01	3.0E+07	8.8E+15	1.4E+00	4.2E+01	2.0E+00	2.2E+14
Mo	1.4E+01	1.7E+07	7.5E+15	1.2E+00	1.7E+01	8.3E-01	1.2E+14
Ag							6.3E+15
Cd							7.7E+15
Sn	7.3E+00	1.1E+07	2.0E+15	3.3E-01	2.4E+00	1.2E-01	1.4E+13
Ba	2.4E+00	5.8E+06	2.0E+14	3.3E-02	7.8E-02	3.7E-03	1.1E+13
Ta							3.4E+12
Pb	5.0E+01	2.2E+07	1.2E+17	2.0E+01	1.0E+03	4.8E+01	
			6.1E+17	1.0E+02	2.1E+03	1.0E+02	3.3E+17

* amount of admixtures in a material of the anode.

Figures 4.8–4.10 present the diagrams with the distributions over the mean velocities, mean energies, and kinetic energies of the ions of chemical elements entering the plasma bunch composition in the experiments with targets and substrates made of various materials. The data are given in the decreasing and increasing orders for the velocities and energies of ions, respectively.

The data of the tables and diagrams give jointly the clear representation of the element composition and power of the ionic and electronic components of plasma bunches in the mentioned experiments. We note that a share of the energy referred to the electronic components of a plasma bunch is slight.

The optical spectra of the discussed experiments contain, besides the lines of a target material, the spectral lines of such chemical elements as H, C, O, Be, Mg, Fe, Ni, Co, Ca, S, P, and others. Moreover, some of the chemical elements registered in the optical measurements (Fe, Zn) compete with the element of a matrix by the amount of emitting ions and the kinetic energy. The mentioned registered chemical elements cannot be referred to the initial target composition, since we used materials of high purity in the experiments, and the measurements of optical spectra were carried out in vacuum at a pressure of residual gases of about 10^{-3} Pa in the chamber of the setup. In our opinion, the absence of a considerable correlation between the composition of the emitting ions of a plasma bunch and the composition of admixtures of the materials of the anode and the substrate (see Table 4.3) testifies to the favor of the hypothesis of a nuclear regeneration of a target substance. This is also indicated by the presence of Fe, being the element with the highest binding energy per nucleon, among the leaders in kinetic energy.

We note the following. In Tables 4.4–4.6, we give the total number of nucleons by separate chemical elements of a plasma bunch in different experiments. By the amount of nucleons among the chemical elements of a plasma bunch in the experiments with a target made of Pb, S occupies a considerable position yielding only to Pb, Cu, and Fe (in this aspect, most characteristic are the results of experiment No. 8197 performed under conditions of the absence of an accumulating screen near the damaged part of a target). In this case, the data of the analysis of the element composition of target explosion products performed with the use of high-precision physical methods show that S plays the role of a “conditionally averaged product” of the detectable part of the regenerated substance of a target in the experiments on Pb.

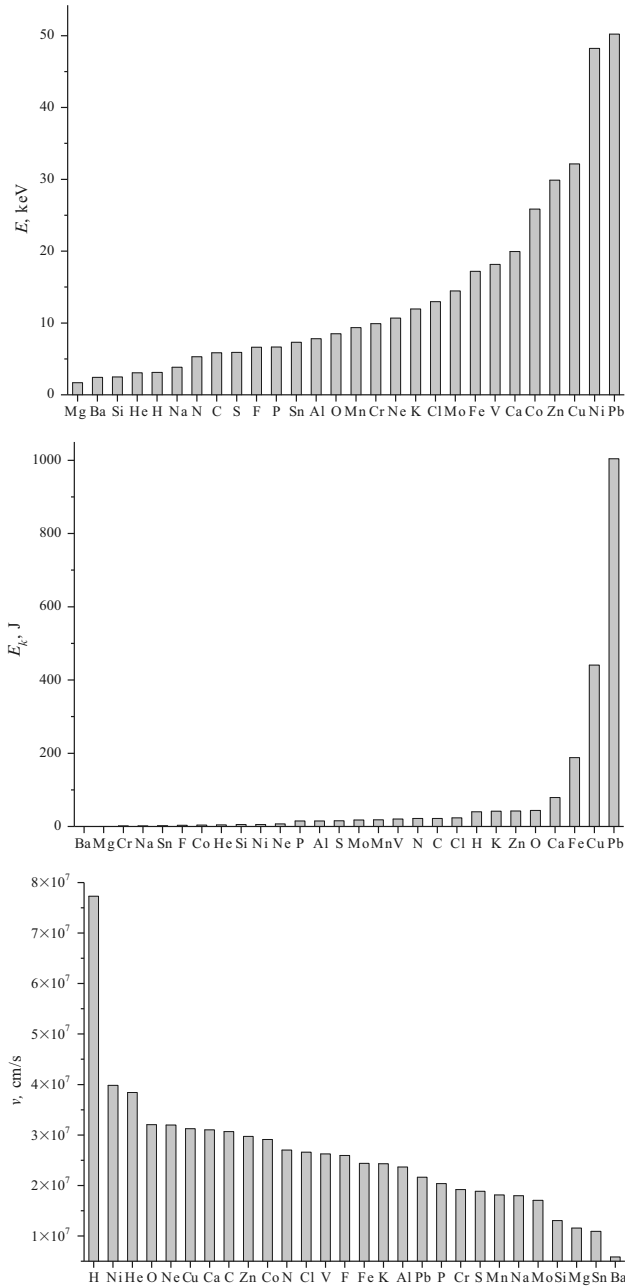


Fig. 4.8. Diagram of the distributions over mean energies, kinetic energies, and mean velocities of the ions of a plasma bunch in the experiment with a Pb target and a Cu substrate.

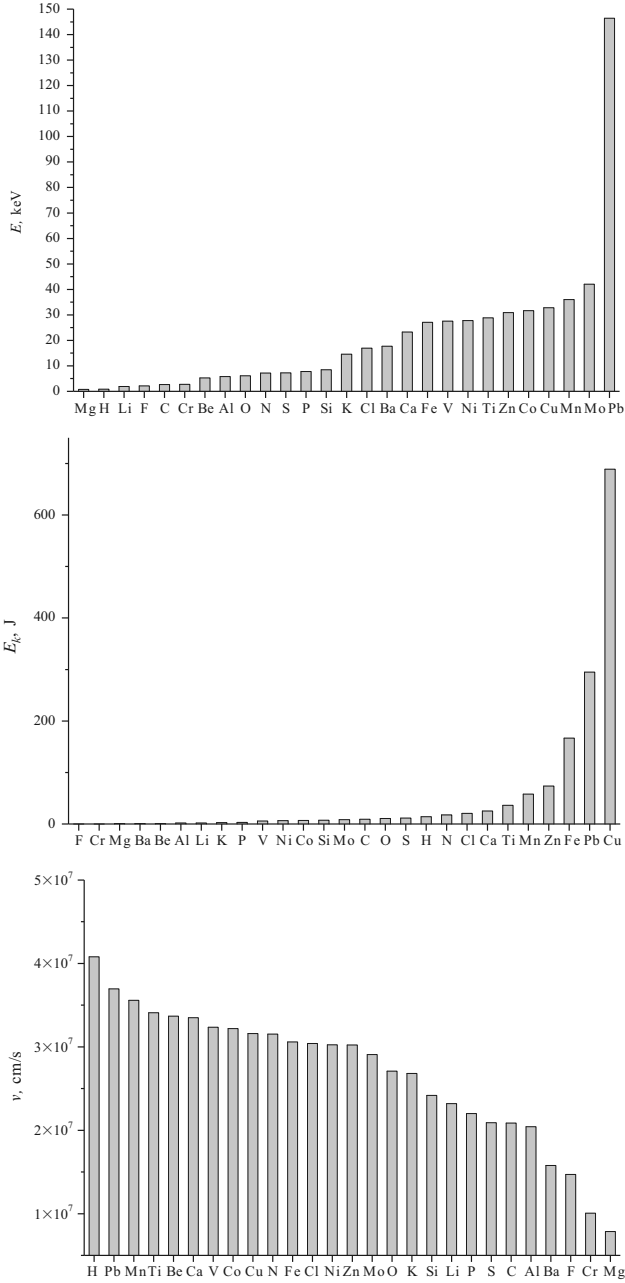


Fig. 4.9. Diagram of the distributions over mean energies, kinetic energies, and mean velocities of the ions of a plasma bunch in the experiment with a substrate and a target made of Cu.

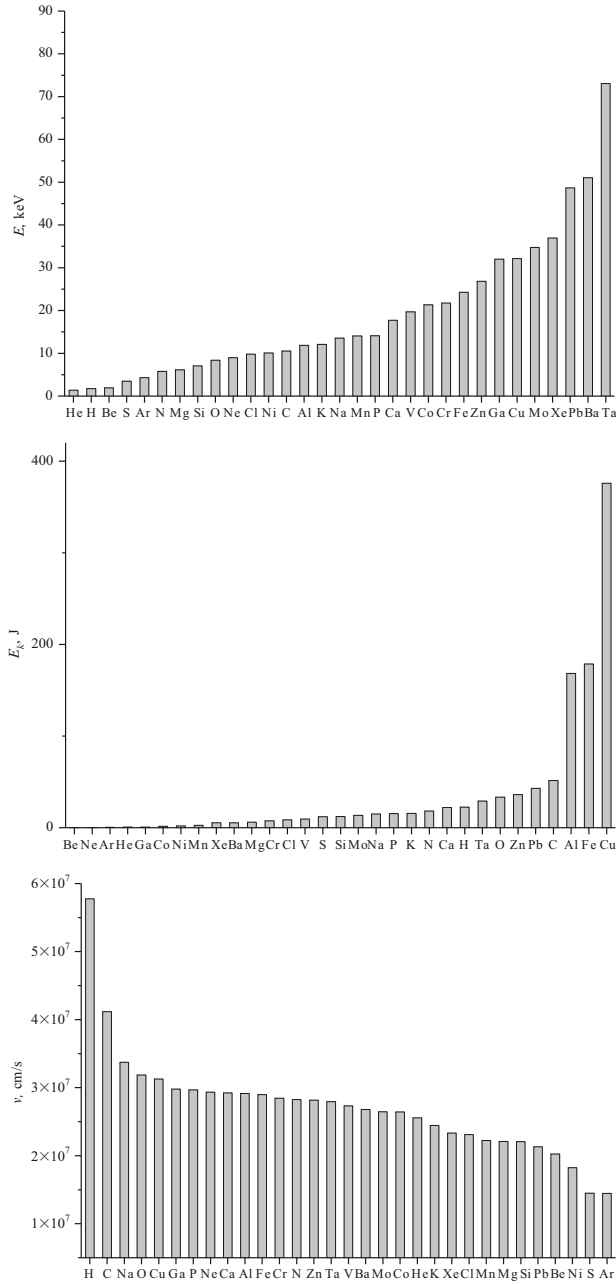


Fig. 4.10. Diagram of the distributions over mean energies, kinetic energies, and mean velocities of the ions of a plasma bunch in the experiment with a Al target and a Cu substrate.

Table 4.3. Correlation coefficients on the quantitative composition of the emitting ions of a plasma bunch and the admixture atoms in the materials of a target and a substrate.

Exper.	Coeff.
6427	0.098
6100	0.022
6466	0.033
6759	0.036
8140	0.015
8134	0.297
8171	0.005
8190	−0.016
8197	0.03

Table 4.4. Nucleon composition of a plasma bunch in various experiments.

	No. 6427 (Pb/Cu)		No. 8140 (Pb/Ta)		No. 8190 (Pb/Al)	
	Number of nucleons	%	Number of nucleons	%	Number of nucleons	%
H	8.10E+16	0.19	6.67E+16	0.41	4.20E+16	0.44
He	3.36E+16	0.08			1.36E+15	0.01
Li						
Be						
C	2.84E+17	0.66	2.36E+17	1.46	3.46E+17	3.67
N	3.63E+17	0.84	1.14E+17	0.70	5.91E+16	0.63
O	5.15E+17	1.19	2.21E+17	1.37	1.84E+17	1.95
F	5.44E+16	0.13				
Ne	8.33E+16	0.19				
Na	6.94E+16	0.16			1.26E+16	0.13
Mg	9.65E+15	0.02			2.12E+16	0.23
Al	3.24E+17	0.75	6.12E+16	0.38	2.73E+17	2.89
Si	3.81E+17	0.88	2.15E+16	0.13	1.45E+16	0.15
P	4.31E+17	1.00	6.00E+16	0.37	5.42E+16	0.57
S	5.23E+17	1.21	2.55E+17	1.57	2.16E+17	2.29
Cl	4.03E+17	0.93	4.29E+16	0.26	1.52E+16	0.16
Ar						

K	8.50E+17	1.96	1.53E+17	0.94	7.05E+16	0.75
Ca	9.93E+17	2.29	2.73E+17	1.68	1.60E+17	1.69
Ti			2.92E+16	0.18	1.41E+17	1.50
V	3.55E+17	0.82	1.12E+17	0.69	1.22E+17	1.29
Cr	5.25E+16	0.12	5.12E+16	0.32		
Mn	6.74E+17	1.56	3.35E+16	0.21	6.27E+16	0.66
Fe	3.81E+18	8.81	8.77E+17	5.41	5.91E+17	6.26
Co	5.16E+16	0.12	6.25E+16	0.39	2.72E+16	0.29
Ni	4.16E+16	0.10	5.41E+15	0.03	3.38E+16	0.36
Cu	5.44E+18	12.57	1.68E+18	10.39	1.10E+18	11.70
Zn	5.76E+17	1.33	1.78E+17	1.10	1.52E+17	1.61
Ga						
Kr						
Mo	7.21E+17	1.67	8.96E+16	0.55	1.34E+17	1.42
Ag			2.73E+16	0.17	8.78E+15	0.09
Cd						
Sn	2.43E+17	0.56	7.94E+16	0.49	3.55E+16	0.38
Xe						
Ba	2.74E+16	0.06	6.34E+16	0.39	2.02E+17	2.14
Ta			4.05E+18	25.01	2.57E+17	2.72
W						
Pb	2.59E+19	59.81	7.36E+18	45.40	5.10E+18	54.00
Total:	4.33E+19	100.00	1.62E+19	100.00	9.44E+18	100.00

Table 4.5. Nucleon composition of a plasma bunch in various experiments.

	No. 8197 (Pb/–)		No. 6466 (Cu/Cu)		No. 6759 (Cu/Cu)	
	Number of nucleons	%	Number of nucleons	%	Number of nucleons	%
H	4.52E+16	0.50	1.49E+17	0.57	1.03E+17	0.59
He	2.10E+15	0.02				
Li					4.26E+16	0.24
Be			1.54E+16	0.06	8.96E+15	0.05
C	2.75E+17	3.02	3.64E+17	1.39	2.52E+17	1.44

Table 4.5. *Continued.*

	No. 8197 (Pb/-)		No. 6466 (Cu/Cu)		No. 6759 (Cu/Cu)	
	Number of nucleons	%	Number of nucleons	%	Number of nucleons	%
N	2.06E+16	0.23	3.21E+17	1.22	2.14E+17	1.22
O	2.26E+17	2.48	2.98E+17	1.13	1.78E+17	1.02
F			7.41E+16	0.28	5.59E+15	0.03
Ne						
Na	1.13E+16	0.12	7.13E+16	0.27		
Mg	2.18E+16	0.24	1.79E+17	0.68	1.13E+17	0.64
Al	1.85E+17	2.03	8.13E+16	0.31	5.42E+16	0.31
Si	3.78E+16	0.42	2.53E+17	0.96	1.47E+17	0.84
P	4.39E+16	0.48	1.23E+17	0.47	7.53E+16	0.43
S	3.62E+17	3.99	3.68E+17	1.40	3.15E+17	1.80
Cl	1.15E+16	0.13	3.35E+17	1.28	2.67E+17	1.53
Ar						
K	6.69E+16	0.74	1.99E+17	0.76	4.42E+16	0.25
Ca	2.08E+17	2.28	4.44E+17	1.69	2.73E+17	1.56
Ti	1.28E+17	1.41	2.23E+17	0.85	3.75E+17	2.15
V	1.58E+17	1.74	6.58E+16	0.25	6.59E+16	0.38
Cr					1.50E+16	0.09
Mn	7.20E+16	0.79	1.75E+17	0.67	5.53E+17	3.17
Fe	9.44E+17	10.38	3.63E+18	13.83	2.15E+18	12.29
Co	5.11E+15	0.06	1.96E+17	0.75	7.84E+16	0.45
Ni	2.30E+16	0.25	2.48E+17	0.95	8.45E+16	0.48
Cu	6.61E+17	7.27	1.25E+19	47.61	8.34E+18	47.68
Zn	3.49E+17	3.84	1.44E+18	5.48	9.72E+17	5.56
Ga						
Kr						
Mo	5.65E+16	0.62	2.98E+17	1.13	1.21E+17	0.69
Ag	1.26E+16	0.14				
Cd	1.27E+16	0.14				
Sn	4.93E+16	0.54				
Xe						
Ba	1.49E+17	1.64	4.21E+17	1.60	3.06E+16	0.18
Ta						
W						
Pb	4.95E+18	54.49	3.78E+18	14.40	2.61E+18	14.93
Total:	9.09E+18	100.00	2.62E+19	100.00	1.75E+19	100.00

Table 4.6. Nucleon composition of a plasma bunch in various experiments.

	No. 8171 (Cu/Cu)		No. 6100 (Al/Cu)		No. 8134 (Al/Ta)	
	Number of nucleons	%	Number of nucleons	%	Number of nucleons	%
H	6.15E+16	0.50	8.09E+16	0.50	5.43E+16	0.94
He			1.12E+16	0.07		
Li						
Be			5.48E+15	0.03		
C	3.59E+17	2.90	3.67E+17	2.28	1.65E+17	2.84
N	4.04E+16	0.33	2.76E+17	1.71	1.99E+16	0.34
O	6.83E+17	5.52	3.97E+17	2.46	9.52E+16	1.64
F	2.85E+15	0.02				
Ne	1.08E+16	0.09	2.72E+15	0.02		
Na	2.05E+16	0.17	1.59E+17	0.99		
Mg	3.79E+16	0.31	1.46E+17	0.90		
Al	7.37E+16	0.60	2.39E+18	14.84	1.55E+18	26.81
Si	7.20E+16	0.58	3.01E+17	1.87	1.47E+16	0.25
P	7.69E+16	0.62	2.10E+17	1.30	1.30E+16	0.23
S	3.62E+17	2.93	6.87E+17	4.26	1.84E+16	0.32
Cl	1.40E+17	1.13	1.93E+17	1.20	5.07E+15	0.09
Ar			2.74E+16	0.17		
K	4.50E+17	3.64	3.16E+17	1.96	5.96E+16	1.03
Ca	1.83E+17	1.48	3.12E+17	1.94	6.13E+16	1.06
Ti	6.05E+16	0.49			2.88E+16	0.50
V	1.08E+17	0.87	1.53E+17	0.95	1.55E+16	0.27
Cr			1.12E+17	0.70	1.60E+16	0.28
Mn	8.52E+16	0.69	6.15E+16	0.38		
Fe	9.64E+17	7.79	2.57E+18	15.94	1.17E+18	20.25
Co	4.64E+16	0.37	2.72E+16	0.17	1.24E+16	0.21
Ni	3.97E+16	0.32	6.75E+16	0.42	2.21E+16	0.38
Cu	2.69E+18	21.76	4.65E+18	28.84	2.02E+17	3.49
Zn	3.82E+17	3.09	5.49E+17	3.41	1.96E+17	3.39
Ga			8.51E+15	0.05		
Kr						
Mo			2.34E+17	1.45		
Ag	1.87E+16	0.15				
Cd						
Sn	6.17E+16	0.50			1.21E+16	0.21
Xe			1.16E+17	0.72		

Table 4.8. Correlation coefficients on the mean energies of the ions of a plasma bunch in various experiments.

Experiment	6427	6100	6466	6759	8140	8134	8171	8190	8197
6427	1	0.476	0.551	0.734	0.66	0.682	0.634	0.698	0.648
6100		1	0.749	0.689	0.696	0.72	0.773	0.739	0.773
6466			1	0.833	0.873	0.894	0.849	0.811	0.851
6759				1	0.977	0.97	0.946	0.971	0.955
8140					1	0.949	0.848	0.959	0.934
8134						1	0.927	0.904	0.959
8171							1	0.837	0.918
8190								1	0.918
8197									1

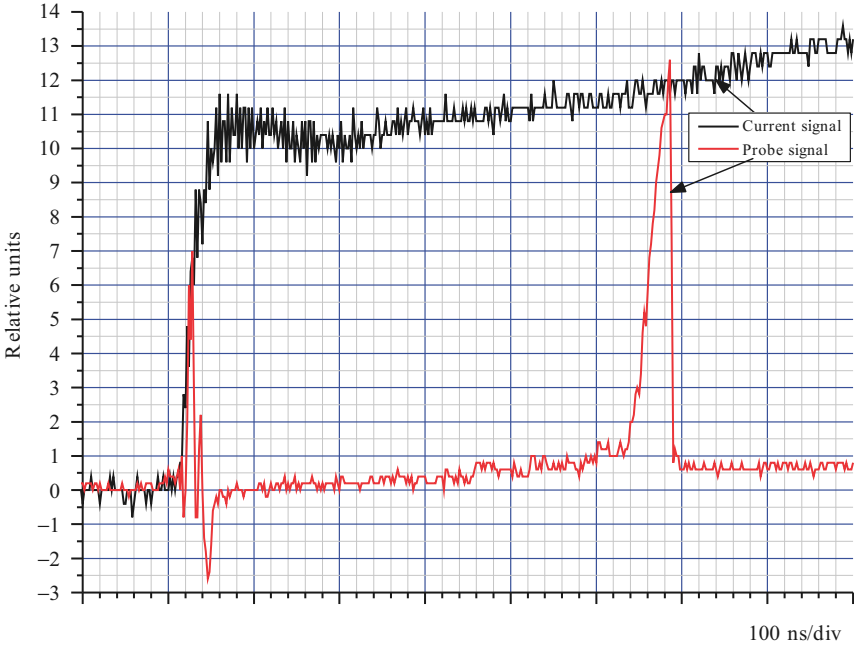


Fig. 4.11. Oscillogram of a signal of the probe registered the arrival time of the plasma front.

probe electrodes was 13 cm; and the time interval from the time moment of the commutation of a current into the diode to the short circuit time of the probe electrodes equaled 574 ns (see the oscillogram). This yields the estimate of the propagation velocity of a plasma front $V \sim 2.26 \times 10^7$ cm/s. As for the highest mean velocities, it is necessary to distinguish hydrogen and

deuterium not separated in optical measurements due to a low resolution of the used equipment, but registered separately with the use of the track method. We may assume that they take a significant share of the released energy. However, this energy is spent to a considerable extent in collisions upon their passage from the collapse zone center through a target material. We connect the observation of the lines of a substrate material in the optical spectra with the action of the emission from HD on the substrate, which causes its erosion. In this case, the substrate atoms are ionized in the plasma bunch and accelerated to the observed velocities by the ions of the regenerated substance of a target.

The most important and significant point, which should attract our attention, is the energy yield of the processes under study. The presented data indicate that almost all the energy yield derived by the results of the analysis of a broadening of the spectral lines of chemical elements entering the composition of a plasma bunch in the optical emission range 300 to 700 nm is represented by the ionic component and is the kinetic energy of the ions of a regenerated substance which are escaping from HD with high velocities. In fact, this fact is one more evidence for the explosive character of the processes running in the vicinity of HD. In this case, the derived results show that the energy yield related to the corpuscular component of a plasma bunch reaches a value of ~ 2.1 kJ (exper. No. 6427) which exceeds significantly the total energy of the electron beam acting strikingly on a target. By the results of systematic electric measurements, the latter is in the range from 100 to 300 J.

4.3. Conclusions

1. The results of the analysis of a broadening of spectral lines of the ions of a plasma bunch indicate the high-energy content of the processes running in the vicinity of HD. The total amount of energy in the corpuscular component of a plasma bunch emitting in the optical range is ~ 2100 J (whereas the energy of the electron beam is ≤ 300 J by the data of electric measurements).
2. The presence of some chemical elements in the composition of a plasma bunch of HD, in particular Fe, Zn, Ca, S, and others, competing by both the kinetic energy and the amount of ions emitting in the optical range with elements of a matrix (a target), as well as the absence of a considerable correlation between the compositions of emitting ions and admixtures of the matrix, testifies to the favor of the hypothesis on the nuclear regeneration of a target material as a result of its shock compression to superhigh densities.

MEASUREMENTS OF X-RAY EMISSION OF HD

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I. V. Skikevich, and B. K. Pryadkin**

The study of the characteristic electromagnetic emission of a plasma presents one of the most important possibilities for the diagnostics of its parameters. In the experiments carried out on the setup IBR-1, there occurs the formation of a dense and hot region, HD, in the anode material under the action of an electron beam. HD reveals all the signs inherent in a strongly compressed plasma: the flash of X-ray emission (XE), high temperatures of electrons and ions, generation of the directed fluxes of accelerated charges, and anomalously large amount of hard X-ray quanta. According to our ideas, these manifestations testify to both the energy concentration extremely high on the terrestrial scale and the approach of the state of matter at the HD to the intrastellar one by the physical properties. In this connection, we get an extraordinary possibility to investigate the superdense matter under conditions of a laboratory experiment. Emission is both the main source of information about the processes running under these conditions and, by means of the withdrawal of energy, a factor defining the dynamics of similar systems. Therefore, the properties of the emission in the state of limiting compression of matter, including the spectral composition, deserve the comprehensive study.

Below, we present the results of the spectral measurements of XE of HDs in a wide range of energies from 10 keV to 3 to 5 MeV.

5.1. Procedure of Measurements

In spectral measurements, we use the method of absorbing filters with variable thickness made of a single material. The spectra of X-ray emission were derived from experimentally measured filtration curves by solving the inverse problem.

Measurements of X-ray emission signals of HD were performed with semiconductor detectors made of GaAs with the sensitive region of $\sim 500 \mu\text{m}$ in thickness and $\sim 0.1 \text{ cm}^2$ in area in the current mode. XE signals were registered with high-speed digital oscillographs “Tektronix” without preliminary

amplification. The measuring block for spectral measurements of XE included six GaAs detectors positioned in a organic glass mandrel in two rows in order to provide a maximally close arrangement of detectors to the center of the measuring channel. Copper filters ($20\text{ }\mu\text{m}$, $50\text{ }\mu\text{m}$, $100\text{ }\mu\text{m}$, 1 mm , 3 mm , and 11 mm in thickness) were fixed on a brass plate with a hole before each of the detectors. The distance between filters and GaAs detectors was $\sim 3\text{ mm}$. The measuring block was located in a steel flange serving, simultaneously, as a seal of the vacuum system and a fixture of the detectors and output connectors. In addition, the flange along with the output window, which was made of Al of $350\text{ }\mu\text{m}$ in thickness and separated the detectors from the vacuum chamber, serve as a screen protecting the detectors from the influence of electromagnetic noises. Signals of the detectors were supplied to a cabin of measurements to the oscillographs by cables positioned in a copper channel. The bias sources for the detectors were arranged directly in the cabin of measurements. The bias voltage supplied on GaAs detectors was 140 to 150 V , which provides both a high ($\sim 3 \times 10^3\text{ V/cm}$) electric field strength in the sensitive region and the attainment of a limit value of the drift velocity of charge carriers, $\sim 2 \times 10^7\text{ cm/s}$, which are necessary to get a high resolution in time ($\sim 10^{-9}\text{ s}$) and to preserve it to be constant upon a change of the voltage on the very detector.

Measurements were carried out in the asimuthal plane of a device IBR-1 at a distance of 23 cm from the axis, with the use of two different schemes of registration of signals (Fig. 5.1):

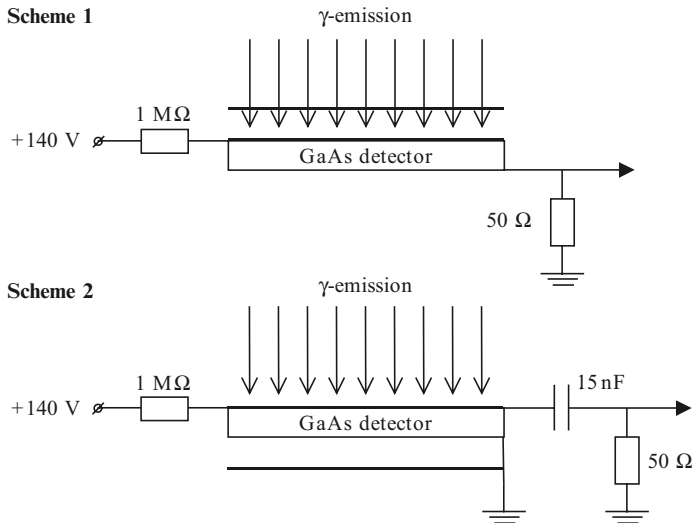


Fig. 5.1. Two different schemes of registration of signals.

Measurements with the use of scheme 1 are more sensitive to the temporal features of signals, because the transmission coefficient for this scheme varies by at most 10% up to the limiting frequency:

$$f_{\text{lim}} \approx \frac{1}{13RC} \left(1 + \frac{R}{R_l} \right), \quad (5.1)$$

where R_l is load resistance equal to 50Ω ; R is resistance of the Ohmic contact of a detector; C is total capacitance of a detector, the oscillograph input, and wires of the circuit.

For a capacitance C of several picofarads, $f_{\text{lim}} \approx 400 \text{ MHz}$ in the most unfavorable case at $R \gg R_l$.

In measurements by scheme 2, the higher amplitudes of signals were registered. In this case, temporal features were smoothed by a separating capacitor.

The energy of X-ray emission absorbed in a detector is spent on the formation of the electric charge collected on the electrodes of a detector with the creation of a voltage pulse on the load resistance. The voltage pulse amplitude on the load resistance and the value of the collected electric charge are related between themselves by a proportional dependence. The main formula for calculations in the approximation of a narrow beam is as follows:

$$\begin{aligned} \alpha \cdot \frac{\epsilon \cdot U(t)}{R_l \cdot S} = & \int_{E_{\min}}^{E_{\max}} E \cdot \Phi(E, t) \cdot \frac{\sigma_A(E)}{\sigma(E)} \left(1 - e^{\sigma(E) \cdot n \cdot \Delta} \right) \\ & \times e^{-\sum_i \sigma_i(E) \cdot n_i \cdot x_i} dE. \end{aligned} \quad (5.2)$$

Here, $S = 0.1 \text{ cm}^2$ is the effective area of a GaAs detector; $R_l = 50 \Omega$; $\Phi(E)$ is the required spectral distribution of XE, U is the voltage pulse amplitude, $\epsilon = 2.67 \cdot 10^{13} \text{ MeV/c}$ is the energy of the formation of an electron-hole pair in GaAs (see Ref. 168), α is a calibration coefficient. The energy of photons absorbed in a detector is proportional to

$$(\sigma_a/\sigma) \left(1 - e^{-\sigma \cdot n \cdot \Delta} \right), \quad (5.3)$$

where σ_a is cross-section of transmission of the XE energy in the detector material, cm^2 ; σ is total cross-section of attenuation of XE in the detector material, cm^2 ; n is density of nuclei of the detector material, cm^{-3} ; Δ is thickness of the sensitive region of a detector, cm .

The subsequent attenuation of the XE spectrum in the anode material surrounding HD, the output window of the setup, and the material of filters is set by the product of exponential functions $e^{-\sigma_i n_i x_i}$, where σ_i , n_i , and

x_i are, respectively, the total cross-section of attenuation of γ -emission, the density of nuclei, and the thickness of each of the mentioned absorbers. On the left and the right (in the integrand) sides of the above equation, the dependence on the filter thickness is represented, respectively, by the voltage pulse amplitude and by one of the exponents describing the attenuation of the X-ray spectrum. A change in the filter thickness varies the absorption of X-ray emission and, respectively, the amplitude of a signal registered by a detector. The dependence of the amplitude of a signal of detectors registered in the spectral measurements on the thickness of the filters used is the so-called filtration function or the apparatus spectrum $U(x)$, where x is the filter thickness. The functional connection between the unknown spectrum of X-ray emission $\Phi(E)$ and the apparatus spectrum $U(x)$ for proportional detectors, if written as

$$\int_{E_{\min}}^{E_{\max}} G(E, x) \cdot \Phi(E) dE = U(x), \quad (5.4)$$

is an integral Fredholm equation of the first kind.

Here, the kernel $G(E, x)$ of the equation is the response function of a detector determined as a result of measurements or by calculations upon the subsequent action of a monoenergetic emission on the system of detectors. The function $G(E, x)$ describes the processes of attenuation of the spectrum of X-ray emission and absorption of the energy of photons of this spectrum in a detector. The calculations of the response function of a detector were performed according to the table data on the cross-sections of interaction of γ -quanta with matter in various physical processes (see Ref. 169). The accuracy of the table data is not worse than 10%.

Equation 5.3 for the energy absorbed in the sensitive region of a detector integrates all the processes of interaction of photons with matter by involving the total cross-sections of attenuation and transmission of energy. This relation does not consider the escape of electrons and photons in separate processes (photoeffect, Compton scattering, and creation of pairs) beyond the limits of a detector. We also did not take into account the ionization of atoms in the sensitive region of a detector by the same particles entering it upon the interaction of the primary emission with materials surrounding the detector, in particular with the materials of filters. Therefore, additionally, to the main contribution defined by Eq. 5.3 to the energy absorbed in detectors, the contributions of the mentioned effects were accounted with the use of the well-known data on cross-sections of the relevant processes (see Refs. 170–173). The rated relations, in which we considered also the process of reflection of charged particles with the use of empiric data (see Ref. 174), are not given here

due to the awkwardness of final mathematical formulas, which would draw away the attention of a reader.

5.2. Results and Discussion

The determination of the unknown function $\Phi(E)$ from the given integral equation is the inverse problem of measurements and is related to the type of ill-posed mathematical problems which can be ambiguous (to have a lot of solutions) or have no solutions at all due to inevitable errors of measurements. Similar problems require special methods of solution. In particular, the Tikhonov variational method of regularization (see Ref. 175) is well known and used. The method is based on the notion of regularizing algorithm as a means of approximate solution. The idea consists in the availability of a certain function possessing the maximum degree of smoothness and preserving the quadratic deviation for the initial equation in the limits of a given accuracy. For this function, we solve the problem of minimization of already the other functional (a smoothing one), in which we introduce a regularizing functional-stabilizer with a coefficient. This coefficient is the regularization parameter defined by the mean quadratic error of measurements. After the finite-difference approximation, the problem of determination of a function minimizing the smoothing functional is reduced to the solution of a system of linear equations. From the functions $\Phi(E)$ derived in the process of solution, one chooses in practice that which has no sharp oscillations and transitions across zero on the least error of the quadratic deviation $U(x)$ by varying the regularization parameter. The choice quality can be verified then by the direct substitution of $\Phi(E)$ in the initial integral equation.

Upon the restoration of the XE spectrum of HD, the studied range of XE from the low-energy edge is limited by an energy of 9 keV. This restriction was related to the absorption of XE in the anode material surrounding the HD. The restriction from the high-energy edge in the X-ray spectrum was not foreseen in the statement of the problem. The hardness of the emitted X-ray spectrum was testified by the very significant amplitudes of signals of a detector beyond the filter of 11 mm in thickness which were registered in some experiments and were equal to $\approx 30\%$ of the amplitudes of signals of the detectors beyond the thinnest filters. In addition, the presence of hard quanta in the MeV range was revealed by the registration of positrons (the creation threshold of electron-positron pairs is 1.022 MeV). In few experiments, we manage to register the activity of silver specimens interpreted by the observed half-life periods of products (active nuclei of the silver isotopes in the amount $\approx 2 \times 10^4$) with the use of photonuclear reactions. The energy threshold of the registered reactions is ≈ 9 MeV. Since no unique deviations were registered by GaAs detectors in these cases, it seems

to be most probable that photons with such an energy are created not in the deceleration processes, but as a result of nuclear transformations in the collapse zone.

Nevertheless, the measurements performed with the use of an analogous collection of Cu filters and a calibrated X-ray “AGFA” film as a detector (naturally, at the same distances from HD) have demonstrated quite inverse. A dose measured by a film is significantly decreased with increase in the filter thickness. The change in the dose was approximately one order of magnitude and more upon the change in the absorbing filter thickness from 20 to 500 μm . The dependence of the dose on the absorbing filter thickness in the “dose in rad – logarithm of the thickness in μm ” coordinates turns out to be linear with the almost identical slope angle in all the performed experiments. With obvious clearness, these results indicate the presence of some inflection in the spectrum in the energy region below 100 keV. This peculiarity was revealed upon the restoration of the spectrum.

While using the method of regularization in solving the problem, we would introduce, first of all, the limitation on the maximum energy in the spectrum. On the maximum energy, for example, of 3 MeV and a step of 10 keV, we would operate with matrices of 300×300 . In solving the problem, the more considerable complexity would be presented by the presence of inflection points and extrema and by their unknown position in the spectrum. For these reasons, the method of regularization in solving the inverse problem, i.e., the restoration of the XE spectrum from experimental data, was not used.

We note the following. From the first measurements, it was found that the increase in the detector signal amplitude beyond a filter of 11 mm in thickness (by several times in certain experiments) leads also to an increase in signals of the detectors beyond the thinnest filters of 20 and 50 μm in thickness, but this increase is less significant. Beyond the remaining filters, the increase in the amplitudes of signals was minimum. This inhomogeneity in the growth of the amplitudes of signals of the detectors beyond various absorbing filters, in our opinion, reflects the existent interrelation of physical processes in the formation of the low-energy (below 100 keV) and high-energy (above 500 keV) regions of the spectrum. Apparently, these parts of the spectrum grow in a correlated way upon the more efficient process of collapse. Just this circumstance is the starting point in the formation of the idea of an XE spectrum as the sum of three components.

In the absence of any restrictions on the maximum energies in a restored spectrum, it seems quite natural to choose the functional energy dependences for separate components of the spectrum in the form of monotonic functions from the well-known elementary spectra (the bremsstrahlung

spectrum of plasma electrons, the black-body emission spectrum, etc.) described by a minimum collection of constants. In this case, any discrete distributions (in the form of separate emission lines) are unacceptable, because they would require a considerably greater collection of constants. Based on these considerations, we chose the same functional dependence on energy in the form of the bremsstrahlung spectrum of plasma electrons for all the components of the spectrum. Such a dependence for the emission of a plasma from the collapse zone is quite obvious and is also acceptable for the electron beam, because the last, under conditions of the pinch-effect reminds most likely a compressed electron relativistic gas with the temperature of electrons of the order of the accelerating voltage supplied to the diode. For higher energies, a similar functional dependence is characteristic of the “temperature” constant $E_3 = e \cdot I/c$, where e is the electron charge, c is the light velocity, and I is a current. We note that the value of this constant falls into the MeV-interval in energy already at the typical current of the electron beam $I \sim 50$ kA. A similar dependence is known as the “induction” spectrum for high-energy protons. However in this case, we consider the choice of such a spectrum for electron-positron pairs and for electrons and positrons, β^- and β^+ , which arise in nuclear decays, are accelerated in the collective fields, and create photons upon the deceleration in the rarefied plasma of a target and, in particular, in the material of filters.

Thus, the *a priori* idea of the spectrum $\Phi(E)$ was realized by the sum of three normed unit spectra represented by exponential functions with weight coefficients

$$(E) = \sum_i C_i \frac{e^{-E/E_i}}{E_1(E_{\min}/E_i)} \cdot \frac{1}{E}, \quad (5.5)$$

where $E_1(Z) = \int_z^\infty (e^{-x}/x) dx$ is the integral exponent function, and $\sum_i C_i = 1$.

Upon the restoration of spectra, we used jointly the experimental data of measurements of the apparatus spectrum by GaAs detectors and an X-ray film, the results of measurements in the low-energy part of the spectrum (in the energy range 10 to 88 keV) with the use of differential filters, and the data of the measurements of doses by dosimeters DK-02 near the detector block.

In the measurements by the Ross method (see Refs. 176, 177), we used five pairs of differential filters made of Cu, Nb, Ag, Mo, Sn, Ta, and Pb foils fully covering the energy range from 8.9 to 88 keV. In the measurements, a detecting element was an X-ray “AGFA” film calibrated from ^{60}Co and ^{241}Am sources in the range of absorbed dose up to 5 rad. The data of the calibration were expanded on the energies in the X-ray spectrum,

which differ from the energies of the calibration sources, with the use of a calculational method given in Ref. 178. The results of measurements by the method of differential filters were used for the normalization of the X-ray spectrum by the intensity at the maximum. In the calculations of doses registered by a DK-02 from the high-energy part of the spectrum (with the energies of photons of above 200 keV), we used the known dependence connecting the exposure dose in air with the density of the incident X-ray emission.

The procedure of restoration of the spectra consists in the substitution of $\Phi(E)$ presented above in the integral equation and in the variation of all coefficients of the accepted approximation of the spectrum up to the coincidence of the calculated apparatus spectrum with the experimental one. In parallel with the use of the restored spectrum $\Phi(E)$, we calculated doses on X-ray films beyond some filters and at the mounting place of dosimeters DK-02. Using the calculated doses on the film, we reproduced its filtration dependence, whose slope was compared to that known from other experiments (see above). The calculated dose at the mounting place of dosimeters was also compared to the measured one. These additional conditions were used for the correction of the coefficients of the approximation of a spectrum and favored the uniqueness of its determination. The accuracy of our measurements is estimated as 10% to 15%, and the used table data have the accuracy close to the above-presented one. The derived results on the X-ray emission spectrum of HD have, in our opinion, the same accuracy.

The derived X-ray emission spectrum of the HD (below, it is compared to the spectra of astrophysical objects) is presented by three components according to the approximation used upon the restoration. Separate components of the spectrum can be interpreted as follows:

- the low-energy component (<100 keV, with a maximum at ≈ 30 keV) corresponds to the spectrum of a hot plasma from the collapse zone;
- the energy region <500 keV includes the contributions of both the bremsstrahlung of the electron beam on a pre-pinch stage and the magnetobremssstrahlung of electrons on the pinch stage;
- the high-energy component (energies >500 keV), for which the most probable mechanism of its formation seems to be the creation of photons in the collapse zone as a result of nuclear transformations.

The energy yield of XE in the spectrum to the full solid angle under isotropic emission is of the order of 1 J with the following distribution of energy over the spectrum: 50%, 70%, and 90%, respectively, in the intervals 10 to 120, 10 to 250, and 10 to 2500 keV.

However as noted above, one of the components of X-ray spectra registered in experiments is the electron beam bremsstrahlung. Upon the deceleration of electrons with a given energy in targets with different atomic numbers, a certain share of the energy of electrons transits into that of bremsstrahlung. According to the data in Ref. 179, this share for electrons with energies of 300 to 500 keV and targets with $Z \approx 30$ is 1.07% to 1.62%. Respectively, at beam currents of 30 to 35 kA and the durations of a pulse of X-ray emission of 12 to 15 ns, the energy yield into bremsstrahlung should be from 1.2 to 4.3 J. But the measurements on a setup IBR-1 showed that the contribution of bremsstrahlung of the electron beam to the registered spectra is less than the rated one. The reason for this fact is unclear else. In addition, since we have discussed, to a certain extent, the high-energy part of the X-ray emission spectrum, it is worth mentioning the observation of the X-ray emission pulses in the recent series of experiments which were shifted by ≈ 70 to 260 ns relative to the usually registered signal at the time moment of the commutation of a current into the diode. The share of such observations is $\approx 3\%$ to 5% of the total number of experiments. In view of the facts that the sampling included only successive experiments and the position of detectors was constant in the process of measurement, we may assume the presence of a possible connection between the registration of these signals and the significant anisotropy of the exit of the high-energy emission from the collapse zone. We also note that the data on the total X-ray energy yield agree with the results of measurements performed with the use of an integral obscure-chamber.

5.3. Comparison of the Spectrum of HD with Those of Compact Astrophysical Objects

According to the directivity of our studies, it seems expedient to carry out the comparison of the measured X-ray spectra of HD to the spectra of some well-known astrophysical objects (the sun, a pulsar from the Crab nebula, quasar 3C 273, the supernova flashed in 1987 [SN1987A]) and to the spectra of short-term flashes of γ -emission in the Universe. Such a comparison is undoubtedly useful in the aspect of the clarification of general regularities in the generation of the emission by HD and by various astrophysical objects. The data on the spectra of astrophysical objects are taken from the literature (see Refs. 180–182). The spectra of astrophysical objects in the X-ray and γ -ranges are separated from their full emission spectra and are presented on the scale convenient for the comparison to the HD spectrum. Below, we give some brief information about these objects.

The solar γ -emission is registered by modern facilities only in the periods of bursts. From 140 registered solar bursts accompanied by the

emission of the measured fluxes of γ -quanta with energies above 300 keV, the energy spectrum containing the registered nuclear lines was measured in 100 bursts. Their appearance is related to the processes of interaction of protons, α -particles, and heavier nuclei accelerated during the bursts with the substance of the solar atmosphere, which results in the excitation of nuclear levels, fission of nuclei, generation of new elements and nuclides. In solar bursts, the emission maximum in the X-ray range has energy of several keV, and the spectrum intensity in the γ -range decreases very monotonously with increase in the energy.

Quasars are extragalactic objects of small angular sizes. The observational data in the whole range of electromagnetic emission are interpreted as follows. Quasars are galactic cores, where the powerful energy release occurs from regions with a characteristic size $\sim 10^{16}$ cm. The emission fluxes in various regions of the spectrum vary in time. The data on the variability of the emission in the X-ray range indicate the extraordinary compactness of the emitting region. There are grounds to assume that the most probable mechanism providing the high luminosity of these objects is connected with the energy release under the falling of a gas on a supermassive “black hole” with $M \sim 10^8$ to $10^9 M_\odot$, where $M_\odot$ is the sun’s mass. Such an interpretation is consistent with both the emission flux variability requiring the source compactness and the presence of the ejections of matter in a definite direction which indicate the long-term stability of a spatial orientation of the source.

According to the generally accepted model, pulsars are rotating neutron stars with a magnetic field on the surface of $\sim 10^{12}$ Gs. A strong anisotropy of the emission is observed. The transformation of the energy to that of nonthermal emission occurs as follows: rotation of a neutron star; appearance of a strong electric field in the vicinity of the neutron star due to unipolar induction; acceleration of particles in the electric field up to relativistic energies; generation of γ -emission upon the motion of ultrarelativistic particles along the distorted magnetic force lines; absorption of γ -quanta in a strong magnetic field; creation of electron-positron pairs; development of plasma instabilities. Young pulsars are located inside the remnants of supernovas (this connection is established for 8 pulsars). For the lifetime of pulsars, the shells around them have already dispersed. About 4% pulsars enter the composition of double systems.

Supernovas are stars, whose brightness increases upon the burst by several stellar magnitudes for several days. The burst energy of such a star is 10^{50} to 10^{51} erg, and the emission power is above 10^{41} erg/s. The luminosity of a supernova can be comparable with that of the whole star system, where it has burst, and can exceed it. About 600 bursts of extragalactic supernovas

are registered. For 100 of them, the curves of brightness (the dependence of brightness on time) and the emission spectra are obtained. On the basis of the observed brightness curves and the optical spectra, supernovas are divided into 2 main types. For supernovas of type I, the absence of hydrogen lines in the optical spectra is typical, and the curves of brightness reveal a noticeable resemblance. Supernovas of type II have hydrogen lines in the spectra, and the curves of brightness are characterized by a significant variety of forms. It is considered that the bursts of supernovas are a result of the dynamical evolution of the star core which begins from the time moment of a violation of the hydrostatic balance in the star and completes by the total dispersion of its substance (supernovas of type I) or by the gravitational collapse of the core (supernovas of type II). In this case, the character of the star's evolution at its last stages is mainly defined by the star mass. The explosive release of energy accompanied by the burst of a supernova leads to the formation of a shock wave propagating to the star surface. The passage of a shock wave causes an increase in the internal energy of a substance acquiring the large expansion velocities, as a result. The expansion of the ejected substance is accompanied by the adiabatic cooling mainly defined by the star radius before the burst. In this connection, the observed brightnesses of supernovas can be obtained at the initial radii comparable with the photosphere radius at the brightness maximum ($\sim 10^4 R_O$, where R_O is the earth's radius). At the considerably lesser initial radii, the existence of an additional source of energy which continuously compensates the adiabatic losses during the expansion of a substance is assumed. Such a source of energy is the decay of a radioactive isotope of Ni in Co and further in Fe. In particular, the observations of the supernova in the year 1987 registered the hard spectrum of X-ray emission which is completely different from those of the known sources of the cosmic emission. By the data in Ref. 181, the character of the X-ray spectrum is a result of the comptonization of the γ -emission upon the decay of ^{56}Co with regard to its mixing in the significant mass of the supernova shell and other effects.

γ -flashes are intense pulse fluxes of γ -quanta with energies from one ten to hundreds of keV propagating in the interstellar space of the galaxy. They were discovered in 1973 as a result of the long-term tracing of the level of the space emission simultaneously from several satellites. At once after their discovery, γ -flashes were observed at most 5×8 times for a year and were considered a rare phenomenon. After the mounting of more sensitive detectors on the interplanetary stations "Venera 11-14", the events are observed every 2 to 3 days. Main characteristics of γ -flashes are: the frequency of their appearance, intensity and temporal structure, energy spectrum and its evolution in the course of a flash, total energy flow, and emission

propagation direction. By the emission intensity, γ -flashes exceed considerably the level of the diffusion background of γ -emission from the whole sky, and their fluxes are higher by several orders than those from the well-known discrete sources. The duration of flashes extends from hundredths of one second to hundreds of seconds. The particular group is composed from very short flashes with $\Delta t \sim 10$ to 100 ms. By various estimations, the size of the emitting region in the last case is ~ 3000 km, which is small on the space scale.

γ -flashes are not reliably identified with astrophysical objects which are visible or known by emission in the other spectral regions. There is also no unambiguous explanation of their origin. The continuous spectra of γ -flashes are satisfactorily described by the relation $(dN/dE) E^{-\alpha} \exp(-E/E_0)$ in most cases, where $\alpha = 0.5$ to 1.5, and the characteristic energy E_0 can be considered as a measure of the emission temperature (see Ref. 183). The value of E_0 can rapidly change in time. Frequently, this occurs in the considerable limits (from 100 to 1000 keV). Such strong spectral nonuniformity revealed in a number of measurements defines a visible temporal structure of flashes. In many cases, there are spectral peculiarities in the form of wide lines in the energy region from 30 to 100 keV. They are assumed to appear in the presence of a strong magnetic field in the source due to the selective absorption of the outgoing emission by the external cooler regions of plasma at the electron cyclotron frequency. The observed frequencies correspond to a magnetic field with $B \approx 3 \cdot 10^{12}$ to 10^{13} Gs.

The nature of bursts causing the generation of the flashes of hard X-ray emission is studied slightly. The available data indicate that the sources of flashes are old neutron stars positioned at distances of ~ 10 to 100 pc from sun's system. Among the possible sources of the energy in neutron stars which can explain, in principle, the energetics of γ -flashes, the following ones are distinguished:

- rotation of a neutron star;
- accretion of matter on a neutron star;
- thermonuclear explosion of the matter enriched by light elements on the surface of a neutron star;
- shear elastic stresses in the solid crust and core of a neutron star;
- nuclear explosion of the nonequilibrium matter of a neutron star enriched by free neutrons and heavy neutron-rich nuclei (see Ref. 184).

The last source is related to a nonequilibrium layer consisting of nuclei overenriched by neutrons and free neutrons which can exist at densities in the interval 10^{10} to 10^{12} g/cm³. The movement of the neutron-rich matter into the region with lesser densities leads to a nuclear explosion. This hypothesis connects the nature of γ -flashes with the very actual problem of nuclear physics, the existence of superheavy nuclei (with $A \geq 300$). Upon the ejection from the surface of a neutron star, these nuclei decay making a significant contribution to the chemical composition of matter in the Universe. It is assumed that such a process must induce the formation of twice “magic” stable nuclei ${}_{114}^{298}A$.

The data of some recent observations indicate the presence of a genetic connection of, at least, a part of γ -flashes with the bursts of far supernovas.

The results of comparison of the X-ray emission spectrum of the HD averaged over ≈ 2500 experiments with the spectra of astrophysical objects are presented in Figs. 5.2–5.5. In Fig. 5.6, we give the results of comparison of the former with the bremsstrahlung spectrum of an electron beam on a massive target.

As a measure of the comparison of spectra, we use the correlation coefficients Z , whose calculation [with regard to the logarithmic scales of the plots of the compared spectra $f(E)$ and $g(E)$] was performed by the following formulas: upon the comparison of spectra by the spectral density of the energy flow,

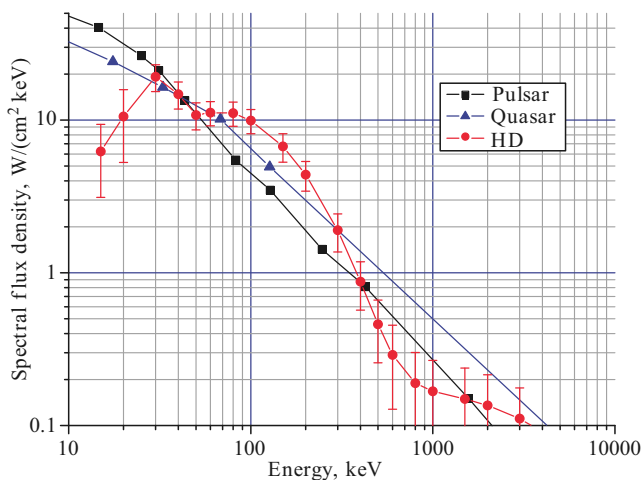


Fig. 5.2. Comparison of the spectrum of HD with spectra of quasar 3C 273 and a pulsar from the Crab nebula by the spectral densities of energy flows.

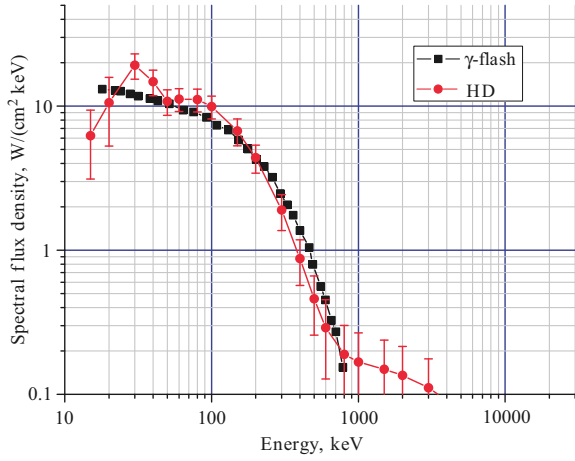


Fig. 5.3. Comparison of the spectrum of HD with that of γ -flashes by the spectral densities of energy flows.

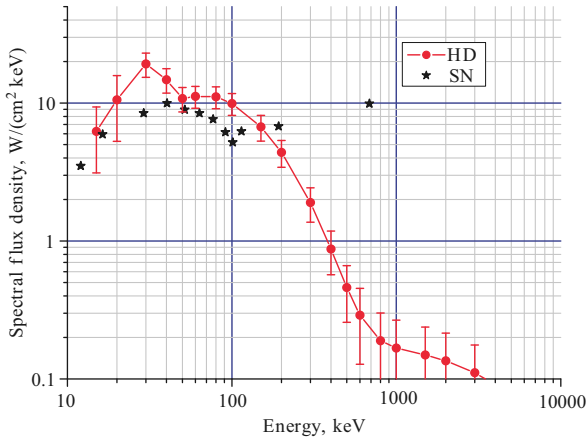


Fig. 5.4. Comparison of the spectra of X-ray and γ -emission of HD with the spectrum of supernova-1987 by the spectral densities of energy flows.

$$\begin{aligned}
 Z &= \int_{E_{\min}}^{E_{\max}} (\lg(E \cdot f(E)) - \langle \lg(E \cdot f(E)) \rangle) \times (\lg(E \cdot g(E)) - \langle \lg(E \cdot g(E)) \rangle) d(\lg E) \\
 &\div \sqrt{\int_{E_{\min}}^{E_{\max}} (\lg(E \cdot f(E)) - \langle \lg(E \cdot f(E)) \rangle)^2 d(\lg E)} \\
 &\div \sqrt{\int_{E_{\min}}^{E_{\max}} (\lg(E \cdot g(E)) - \langle \lg(E \cdot g(E)) \rangle)^2 d(\lg E)},
 \end{aligned}$$

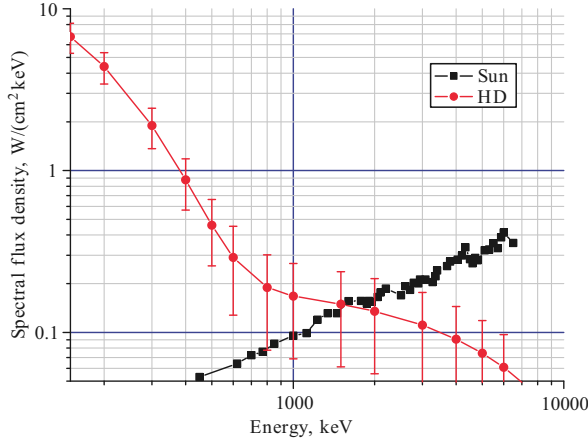


Fig. 5.5. Comparison of the spectra of X-ray and γ -emission of HD with the spectrum of the sun by the spectral densities of energy flows.

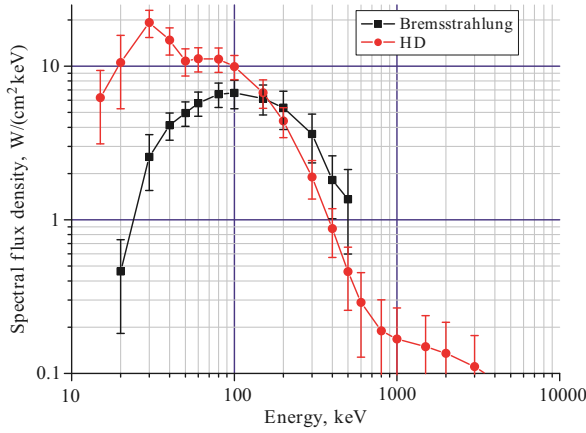


Fig. 5.6. Comparison of the spectra of HD with the bremsstrahlung spectrum of an electron beam by the spectral densities of energy flows.

where

$$\langle \lg(E \cdot f(E)) \rangle = \frac{\int_{E_{\min}}^{E_{\max}} \lg(E \cdot f(E)) d(\lg E)}{\int_{E_{\min}}^{E_{\max}} d(\lg E)},$$

$$\langle \lg(E \cdot g(E)) \rangle = \frac{\int_{E_{\min}}^{E_{\max}} \lg(E \cdot g(E)) d(\lg E)}{\int_{E_{\min}}^{E_{\max}} d(\lg E)}.$$

Table 5.1. Correlation coefficients of the spectral densities of the energy flows of astrophysical objects, bremsstrahlung of an electron beam, and HDs.

Compared objects	Energy interval, keV	Correlation
“Hot dot”, Pulsar	10 to 4000	0.916
“Hot dot”, Quasar	10 to 4000	0.937
“Hot dot”, Sun	200 to 5000	−0.955
“Hot dot”, γ -flash	20 to 800	0.989
“Hot dot”, SN1987A	10 to 700	−0.231
“Hot dot”, Bremsstrahlung	10 to 500	0.243

The data on values of the correlation coefficients and on the energy ranges used in the comparison of X-ray emission spectra are given in Table 5.1.

Thus, the comparison of the spectra of X-ray emission by the spectral densities of energy flows shows a similarity of the spectrum of HD and the spectra of astrophysical objects (a quasar, a pulsar, a γ -flash) and a difference from other spectra (those of SN1987A, the sun, and bremsstrahlung) with high correlation coefficients.

5.4. Conclusions

1. The emission parameters of HD in the energy range from 10 keV to 3 to 5 MeV are very close to those of the emission of nonstationary astrophysical objects (quasars, pulsars, and γ -flashes) which are characterized by the compactness of the emitting region (on the space scale) and the presence of significant magnetic fields.
2. The difference in the spectra of HDs and SN1987A is mainly observed in the high-energy part, where a considerable deficit of photons is observed in the spectrum of HD.
3. The spectrum of HD is essentially different from that of the sun.
4. The spectrum of HD differs from a bremsstrahlung one by the enhanced yield of photons in the low-energy part and by the presence of a long high-energy tail.

6

REGISTRATION OF FAST PARTICLES FROM THE TARGET EXPLOSION

A. A. Gurin and A. S. Adamenko

6.1. Introduction

Together with other diagnostic tools, track dielectric detectors turn out to be a very convenient instrument giving the important information on the intensity and the composition of a corpuscular emission represented by the plasma fluxes flowing out the collapse region of a damaged target. In view of the smallness of this region, we will say about the HD of a pulse discharge, as is customary in the researches of discharges of the “plasma spark” or “plasma dot” type in vacuum diodes (Refs. 185–187, 191). The other related object is a contractive pinch state of the “plasma focus” type (Refs. 188, 189).

In the discharges realized in the framework of the project “Luch”, HD appeared not in a gas-plasma medium, like in plasma diodes or plasma foci, but on the axis of a solid target, usually a metallic one. It is worth noting the well-known approaches to the realization of plasma flashes on the basis of the concentration of energy flows on a microtarget with the help of a relativistic electron beam (Refs. 190, 192) or a focused laser beam. Refs. 185–190 are considered pioneer works and surveys (Refs. 191, 192). The scenario of discharges realized in the experiments of the project “Luch” is close to the schemes used in Refs. 185–189 by the role-playing by processes of self-organization of a collapse and uses no “apparatus” cumulation of beams on a target with the purpose of its compression and heating like that in Ref. 190. The working principle for the experimental setups of the project “Luch” is based on the fulfillment of the initial and boundary conditions for the formation of the electron beam which are set in the “know-how” of the forming device and units of a diode. Track detectors give the additional proofs of that the collapse of targets turns out to be so profound that, in this case, we reach the level of nuclear-physical phenomena possessing the anomalous, earlier unknown nature.

With the help of track detectors installed in an ionic obscure-chamber, we have succeeded to get a quite original split image of HD in the

form of several similar spots: well-pronounced parallel stripes, namely, the track aggregates which are oriented, as a rule, along the axis of a discharge. Every of the stripes begins on the detector in the region of the obscure-chamber near its axis, where a stripe is formed by the particles which are most fast for it. This image of HD testifies to specific, not completely understandable, properties of the realized profound collapse of a target. It is clear that the magnetic field, which is present during the dispersion of products from the damaged target, can serve as an efficient analyzer of the beams of released fast particles. As accumulating screens for the corpuscular emission, track detectors do not give the dynamic pattern of a collapse, but allow one to analyze the composition of the beams of fast particles which are separated by the obscure-chamber and bear the split image of HD.

The first result given by the application of track detectors consists in that the main component of the corpuscular emission registered by the tracks of the thin beams of fast ions is hydrogen ions. The track analysis shows that hundreds of track images of HD derived on detectors in obscure-chambers, which were accumulated during two last years from the beginning of investigations of tracks, were formed by nuclei of hydrogen with paths in a detector CR-39 from several to tenths of microns. In this case, separate stripes in the multistripe aggregates present on the obscurograms after successful shots are formed by one-type particles. The question arises: How do several separate stripes which have clear boundaries and are formed by the same particles during one pulse discharge appear? We can make two assumptions: the stripes are formed by particles somewhat different by their ionization properties or the identical particles are emitted as a pulse at different times in the scope of one discharge on the setup. The track analysis reveals a very interesting fact in this pattern: it turns out that the tracks in the stripe parts comparable by the number and density of tracks belong to two sorts of hydrogen ions and can be interpreted as the tracks of protons and deuterons. Moreover, deuterons are represented by two stripes in three-stripe “triplets”. These results are derived both with the help of obscurograms and on the basis of the analysis of the traces of split beams on detectors located in the external magnetic analyzer.

The direct count of tracks of a thin beam on detectors in the magnetic analyzer allows one to determine the energy spectrum of ions and to estimate the total number of track-forming particles emitted by HD to be of the order of 10^{14} (in the spectrum region represented by the “specific” energy $E^* = (M/Z^2)E > 100 \text{ keV}$, where M and Z are the mass and charge numbers of ions). The last measurements of tracks performed with the help of a Thomson mass-spectrometer, in which thin beams are analyzed in external constant magnetic and electric fields, confirm the dominance of hydrogen

ions in the corpuscular emission of HD. In this case, we can clarify the peculiarities of the presence of the deuterium component in the total emission spectrum. The $M/Z = 2$ parabola, which is a locus of deuterons on detectors in a Thomson analyzer, is always present together with the $M/Z = 1$ parabola formed exclusively by protons. The fact that almost 100% of the track array for the $M/Z = 2$ parabola belong to deuterons is confirmed by the track analysis, including the comparison with the rated data on trajectories for our analyzer. The proton parabola dominates always by the total number of tracks. It is monotonically filled with decrease in energy, but the segments of the $M/Z = 1$ and $M/Z = 2$ parabolas corresponding to the energies $E^* \approx 0.5$ to 0.75 MeV turn out to be filled by tracks with the same density. *That is, protons and deuterons are represented in the proportion 1:1 in a certain energy range.* The total number of deuterons is of the order of 1% of the number of protons with the energy $E^* = E > 0.3$ MeV registered by the used Thomson analyzer.

Here, we give also the data of the first measurements of tracks in the scope of the $M/Z = 3$ parabola on the track detector corresponding to the energy $E^* \approx 1$ MeV, for which the proton and deuterium parabolas are usually empty. The corresponding stripe at this energy is seen only by using a microscope. It is filled by tracks of two types with a density which is significantly less than the absolute one after the 4-h etching, i.e., we are far from the overlapping of craters. The result of track analysis is more striking as compared to the case of deuterons: tracks belong to the nuclear families of hydrogen and helium, and thus *the ions that induced them are unambiguously identified as the nuclei of tritium and ^6He .* The numbers of tritons and the nuclei of ^6He are approximately related as 1:10, but it is else difficult to estimate the total number of these nuclei for the given construction of a Thomson analyzer. Apparently, this number is less by two orders than the number of registered deuterons.

Thus, track detectors help to establish the anomaly of the composition of light ions emitted by HD: the enhanced content of deuterium and the presence of *the nuclei of tritium and ^6He of the artificial origin* in the experiments with a metallic target not subjected to any preliminary ion implantation. These reliably established facts together with other experimental results testify to the running of anomalous nuclear processes in a collapsing target on setups of the Electrodynamics Laboratory "Proton-21".

The application of track detectors as a measuring tool is connected with the design and use of the corresponding efficient method. The necessity of such a method is defined also by the circumstance that the analysis of flows of fast ions with velocities $>10^8$ cm/s with the help of analyzers registering the electric current of fast ions *in situ*, in the beam separated

on the front of the expanding plasma of a damaged target, turns out a difficult problem. A more difficult problem is the time-of-flight analysis of the individual signals from track-forming particles moving with velocities of the order of 10^9 cm/s on the flight base of about 10 cm in the scope of a discharge chamber. Here, the main obstacle is strong electromagnetic interferences induced by a discharge upon taking off the electric signals in the time interval of about 20 ns including the diode breakdown. The measurements of tracks cannot replace the time-of-flight analysis of the signals from ions, but help to unambiguously identify the tracks of fast ions registered in these experiments.

This purpose is reached on the basis of the Somogyi's method of "the asymptote of track diameter squares" (see Ref. 193). Just the last method (named the TDS method in what follows) allows one to make track detectors sensitive measuring instruments additionally to their role of the indicators of the fluxes of fast particles successfully played by them. Apparently, it is the simplest accurate means to analyze the short tracks which are observed as well-measurable objects with diameters of several microns, being already overetched. The TDS method is directly adapted to the determination of the integral characteristics of tracks (the total length R , limiting depth L , and mean relative rate of etching V) which are sufficient for the identification of track-forming particles in the presence of the corresponding calibrating data: RV - or RL - "loci". We know no cases of the application of the TDS method, except for examples presented in Ref. 193, because of its laboriousness and the necessity to study deeply overetched tracks.

The other approach developed in the methodically irreproachable cycle of works (Refs. 194–199) consists in the measurement of the evolution of the depths of tracks on the breaks of detectors. Local values of the etching rate for the tracks of light ions measured in the mentioned works allow one to construct, in principle, any calibrating loci of the integral track characteristics convenient for the identification of particles. However, the region of low energies less than 1 MeV/nucleon remains poorly studied.

We will demonstrate the efficiency of the used procedure by the example of the construction of an RV -locus for α -particles with paths from 3 to $14\ \mu\text{m}$. Though the α -emission is accessible for the calibration, the construction of an experimental locus by the most exact results of direct measurements of the etching rate of the tracks of α -particles (see, e.g., Ref. 196) for values of $R < 12\ \mu\text{m}$ is impossible, as will be shown below. The increase in the mass and the charge of nuclei leads to a growth of not only the etching rate of tracks, but of the error of measurements of V , especially in the neighborhood of its maximum which is positioned just in the region of small paths. The use of the TDS method leads to the conclusion

that, from the methodological viewpoint, it is expedient to base the track analysis of the beams of particles heavier than the nuclei of hydrogen and helium on loci in the coordinates (R, L) , where L is the limiting depth of a track determined directly and most exactly in the process of long-term etching by the TDS method.

The scarcity of data for ions with paths of several microns, which would be suitable for the practical use with the purpose to identify the main component of the corpuscular emission of “hot dots”, is explained by that the corresponding “subnuclear” region of parameters was not of interest for the high-energy physics, where the methods of track registration were mainly developed. Measurements of short tracks are connected with certain difficulties, because the accurate measurements of the fragments of a track cannot be carried out even on the breaks of detectors. This is especially justified for nuclei of hydrogen, whose ionizing ability is low. They do not create deep and slanting tracks, so that the single parameter available for accurate measurements is the diameter of an overetched track on the surface of a detector. We mention work (see Ref. 194) unique in the whole literature on tracks which contains the sufficiently exact data needed for the construction of an experimental R - V -locus of low-energy protons in the region of paths from 3 to 13 μm in the form of 5 points in the coordinates (R, V) . No detailed data for the tracks of low-energy deuterons are available.

The problem of the deficiency of calibrations is solved by the construction of theoretical curves $R(V)$ for all isotopes of hydrogen by representing the etching rate of tracks in terms of the rate of limited energy losses of particles due to ionization, whose dependence on the nuclear parameters is given by the known semiphenomenological model calibrated by the most accurate data (see Ref. 194) for protons. This commonly accepted approach does not lead, however, to the final result for the hydrogen family, because, using the TDS method, we have to construct calibrating loci within the same method. All our track measurements performed by the TDS method and the attempts to realize the calibration-involved constructions of VR -loci with the use of the secondary beams of protons and deuterons derived by the “forward” scattering (by 30°) of 2.7-MeV beams through a Cu foil on a tandem-accelerator at the Institute of Nuclear Research of the NASU in Kiev, show that the experimental values of V_{ex} are shifted downward by 10 to 20% relative to the values on loci calibrated by the data (see Ref. 194) in the region of the corresponding paths.

The reason for this disagreement should be sought in the following. Upon microscopic measurements, we observe not a real track, but a diminished contrast image, whose diameter D_{ex} is less than the real D (see Refs. 196, 199). This optical effect depends on the geometry of a track,

i.e., on the sort of particles. For protons, the difference $D - D_{ex}$ is minimum. At the beginning of the etching of the tracks of protons with an energy ≈ 1 MeV, its value is approximately equal to $1 \mu\text{m}$ (see Ref. 199). To the drawbacks of the TDS method, we refer the fact that the presence of a constant error upon the measurement of D leads to an increase in the error of the experimental values of R and V proportionally to the very value of D . This effect is especially pronounced upon the measurements of small etching pits such as overetched hydrogen tracks. It is evident that this effect is additionally increased with a growth of the difference $D - D_{ex}$ in the process of etching.

Because the results of calibrations of the loci constructed within the TDS method for hydrogen depend on the etching duration or the thickness of the etched layer of a detector and on the energy of particles or values of the corresponding paths, we restrict ourselves only by the estimations of values of a displacement and a splitting of experimental loci. We emphasize that the main advantage of the TDS method allowing one to distinguish the tracks of protons and deuterons is preserved under real limitations. Moreover, the relative intensification of the effect of the difference $D - D_{ex}$ for deuterons as compared to protons helps us to explain some features of the splitting of experimental loci.

On the basis of such a semiquantitative analysis, we conclude that the beams of particles bearing the multistripe images of HD at the ion obscure-chamber and presenting the two-stripe loci after the processing of detectors within the TDS method are protons and deuterons.

We do not discuss the possible mechanisms of generation of the anomalous amount of deuterons, tritium, and ^6He . Here, our purpose consists in the description of the experimental and methodical aspects.

Besides the category of nuclear tracks which are formed by dense jets of fast ions of hydrogen, whose identification was discussed above, we observed the other type of tracks after some shots. By the etching rate, those tracks are close or identical to those of α -particles familiar by observations of the natural α -background. These track aggregates appear in the detectors shadowed from the direct plasma irradiation. Moreover, the orientation of tracks is not connected with the projection on the direction to the discharge center.

As distinct from fast hydrogen fluxes creating the images in obscure-chambers in a foreseen manner and under controlled conditions, tracks of the second category arise seldom and in an unpredictable way. Similar aggregates of nuclear tracks are divided, in turn, into two types:

1. completely chaotic tracks qualitatively perceptible as a sharp increase in the α -background, the main distinction from which consists in the

limitation by the duration of the experiment and the locality of a manifestation;

2. centered track families or clusters, the tracks in which are clearly oriented to the own local center of dispersion.

Below, we present the examples of the position of detectors, on which we observed chaotic aggregates of α -tracks, and the images of track clusters of different powers. We focus our attention on a “giant” cluster registered on the control detector at a magnetic analyzer. The analysis of the dispersion directivity and the simulation of the cluster testify to the synchronism of the act of emission of particles which is analogous to the burst of a hypothetical superheavy nucleus.

The content of this chapter is the following. In Sec. 6.2., we give the general characteristic of the irradiation of detectors exposed to the flux of the expanding plasma of a damaged target. Section 6.3. is devoted to the method of track analysis, and Sec. 6.4. contains the illustrations of the obscurograms of HDs. Here, we give also the examples of the constructions of loci of the multistriple images of HDs on detectors at obscure-chambers and analyze the tracks on detectors in the magnetic analyzer. The images of track clusters and the characteristic of nuclear tracks are presented in Sec. 6.5., and Sec. 6.7. contains a brief discussion and concluding remarks.

6.2. Characteristics of the Corpuscular Emission of an HD

In our investigations, we used detectors CR-39 TASTRAK produced by the Track Analysis Systems firm (TASL, Bristol). We took the standard conditions of etching: etching reagent – NaOH, molarity – 6.25 M, etching temperature – $70 \pm 0.5^\circ\text{C}$, and etching rate was varied in the limits of 1.5 to $1.7 \mu\text{m/h}$. The track detectors, being placed at the distance of 3 to 13 cm from the discharge center and directly irradiated by a plasma front, show the full filling by tracks of various sizes. The results of the irradiation by plasma are shown in Figs. 6.1 and 6.2.

The main background in Fig. 6.1A consists of weak plasma-induced tracks with the absolute packing which are overetched after etching for 5 to 6 h, so that the diameters of tracks cannot be measured, and the total depth of the etched layer is 2 to $3 \mu\text{m}$. Against this background, we see the boundaries of the jets of ions that have created relatively large deeper tracks.

Upon the panoramic incidence of a plasma flash on detectors surrounding the discharge axis as a ring, we observed a significant inhomogeneity of the dispersion of particles creating these “subnuclear” tracks. Moreover, the density of the latter becomes absolute only in the scope of about 120° to 180° of the total azimuth angle in the equatorial plane of

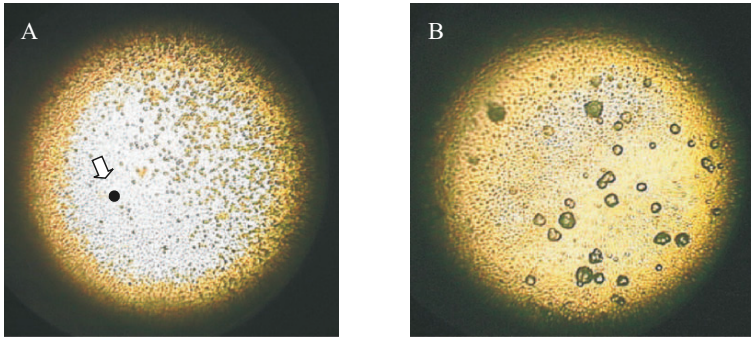


Fig. 6.1. Ordinary picture in the $\varnothing 400\ \mu\text{m}$ field of a microscope (magnification 120) with the track filling of detectors exposed openly at the distance from 3 to 10 cm from the discharge center after etching for 5 to 6 h (A); tracks of especially large plasma particles which are the agglomerates of irregular hemispheres with diameters up to $100\ \mu\text{m}$ (B).

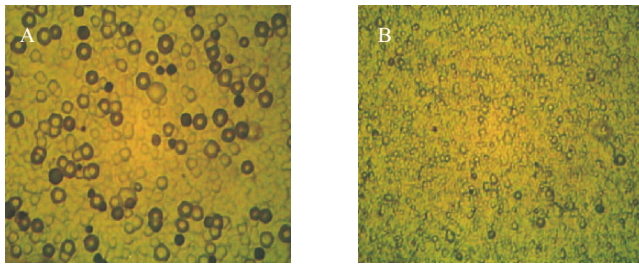


Fig. 6.2. Pattern of tracks after the 6-h etching of detectors exposed in the presence of the intense fast component of the corpuscular emission of HD: normal incidence of particles (A); the angle of incidence is equal to 60° (B).

the diode. Namely these fluxes induce the contrast images at ion obscure-chambers and are identified, as we show in Sec. 6.4., as nuclei of hydrogen, namely, protons and deuterons with paths in a detector from 3 to $14\ \mu\text{m}$. In some experiments on the detectors in obscure-chambers, we registered the stripes formed by larger tracks. The subsequent analysis of similar tracks observed in the Thomson analyzer showed that these tracks belong to particles heavier than the nuclei of hydrogen. In experiments with copper and other metallic targets, the paths of protons and deuterons are usually close to $10\ \mu\text{m}$. But, in separate experiments, we registered intense jets of protons with paths of $20\ \mu\text{m}$ and more (the highest registered energy of protons or, possibly, deuterons is $2.7\ \text{MeV}$ with a path of almost $100\ \mu\text{m}$). The most productive sources of the intense hydrogen emission are the targets produced of polyethylene or heavy metals. In particular in experiments with Pb targets, we registered most stably the fluxes of nuclei of hydrogen with the absolute

packing already at the sixth hour of the etching on detectors remote by 3 to 4 cm from HD. For such targets, we established the splitting of the loci of track-forming particles and the equal representation of protons and deuterons in the analyzed groups of tracks.

An additional information about the character of the process responsible for the generation of high-energy particles was derived in the experiments with hollow targets. The aim of these experiments was the registration of explosion-produced particles upon the action on a target, in which its axisymmetric central part where the nuclear processes are running at HD is absent. However, we do not exclude the possibility for the nuclear processes to run in the hollow on the axis of the target at the expense of the acceleration of ions from the collapsing shell. It is obvious that, in this case, the volume where the nuclear regeneration of a substance occurs, the number of high-energy particles, and the upper bound of their energy spectrum should be considerably lesser. For the mentioned experiments, we chose the Pb targets, whose external surface had the same sizes as those in the experiments, in which we registered the most intense flux of protons with energies of about 1 MeV. The images derived on detectors have confirmed our assumptions. With the use of hollow targets, the total number of track-forming particles on the detectors was lesser by several times. In this case, the maximum energy of registered particles was at most 100 keV per nucleon, and there were no particles with higher energies. The analysis of such weak tracks is a complicated problem, and therefore we cannot answer the question about which nuclei correspond to the registered particles.

Similar experiments are a reliable evidence for that the high-energy particles are generated only in the case of the formation of a collapse at the center of a solid target. Upon the violation of the continuity of the central part of a target, the final stage of the process of collapse is not realized at all or the number and energy of particles participating in this part of the collapse are lesser, which leads to a decrease of the yield of fast secondary particles and their mean and maximum energies.

In the experiments with detectors irradiated openly by a plasma flux of normal incidence, we observed, as a rule, their full filling by overetched tracks. So, it is difficult to separate the tracks with longer paths on this background.

Sometimes, we observed nuclear α -tracks similar to the background on detectors openly exposed (in Fig. 6.1A the arrow indicates a conic track of $20\ \mu\text{m}$ in diameter). In Fig. 6.1B we can see large tracks of the type of etching pits with complicated form with a diameter of $50\ \mu\text{m}$ and more which appear at once as overetched objects and, possibly, belong to fast clusters or microparticles. Such tracks are observed mainly in the near-anode region

of a discharge upon the panoramic registration of the plasma fluxes of HD. But they are registered very seldom inside obscure-chambers. We have else no models and calibrating means for the identification of these particles.

By referring to the detectors exposed openly, we cannot but present the data of a series of experiments, in which the detectors were mounted at a distance of several centimeters from HD and were covered with a fluoroplastic film of $10\text{ }\mu\text{m}$ in thickness. This filter allowed us to get rid of the full filling of a detector by relatively weak tracks. In these experiments, detectors were filled, in dependence on the target configuration, by small long-path hydrogen tracks or scanty large overetched tracks. The study of these large tracks (their diameters were about 12 to $15\text{ }\mu\text{m}$ after the 6-h etching) derived in a Thomson analyzer (we describe them below in detail) show that such tracks are created by the particles with Z equal to at least 5–6 and with paths of about $6\text{ }\mu\text{m}$. By assuming that these tracks are formed by the nuclei of carbon, the energy of one nucleus should be equal to 4.7 MeV (about 400 keV per nucleon). We should not forget that a nucleus has the mentioned energy after the passage of the fluoroplastic filter, where it loses about 8 MeV.

As for the mentioned series, a special attention should be given to the result of experiment No. 6757. After the 6-h etching, we revealed several groups of tracks on a detector, whose diameters were about $25\text{ }\mu\text{m}$ (see the photo). The groups of tracks are concentrated on the detector in circles up to 0.8 mm in diameter. In this case, the tracks with the maximum diameter are concentrated at the center of a circle and become degraded to its margin. Similar tracks belong, apparently, to nuclei heavier than those of carbon or oxygen with residual energy (after the passage of a $10\text{-}\mu\text{m}$ fluoroplastic film) of about 400 to 500 keV per nucleon. Unfortunately, we did not measure the thickness of an etched layer of the detector in this experiment and were not able to performed the exact track analysis.

The registration of tracks on open detectors does not allow us to estimate the total number of track-forming ions emitted by HD along all the directions, because the absolute density of tracks ($>10^8/\text{cm}^2$) is reached at all distances from the discharge center which are available in our measurements (up to 36 cm). Therefore, the number of particles estimated to be 10^{12} for the absolute packing of tracks turns out to be clearly underestimated as compared to the total number of track-forming particles emitted by HD along all directions. The lower energy threshold for the formation of tracks is also unknown. There are indications in the literature that this threshold for CR-39 is 10 keV/nucleon, which corresponds to a velocity of 10^8 cm/s (we have verified that ions of carbon and oxygen with an energy of 40 keV produce no etchable tracks on our detectors). The estimation of the total

energy content of the fast component of the corpuscular emission of HD is additionally complicated by that it depends on the possible contribution of ions heavier than nuclei of hydrogen, but relatively slow. The available data allow us to assert that if the dense fluxes of such ions are registered by open detectors, this occurs only on the track formation boundary. It follows from the data on the registration of the electrical current from the ion component of the expanding plasma of a damaged target that the main ion signal of the products heavier than hydrogen is transferred with a velocity not exceeding 10^8 cm/s. Any exact track measurements are impossible in this region of parameters.

Atomic ions of hydrogen, being the main component of the emission from HD registered by detectors, are also represented in the low-energy range adjacent to the track formation boundary. The domination of the background of nuclei of hydrogen with paths up to $2\mu\text{m}$ and heavy ions with submicronic paths is easily verified by the comparison of the results of the irradiation of open detectors, whose planes are normal to or are directed at an acute angle to the direction to the discharge center (see Fig. 6.2 A and B).

After the 6-h etching, the relatively deep underetched tracks of nuclei of hydrogen dominate on the detector. The depth of these tracks is 2 to $3\mu\text{m}$, and their diameters are from 4 to $8\mu\text{m}$ depending on the energy of particles in the case of the normal incidence of a dense jet of the main component of the nuclear emission (Fig. 6.2A).

Let the angle of incidence θ be 60° . According to the main condition for the observability of tracks $V \cos \theta > 1$, where V is the relative etching rate, the hydrogen atomic component should not be observed in the general case: for nuclei of hydrogen, according to the known data, the condition $1 < V < 2$ is satisfied. Figure 6.2B demonstrates the full filling by small overetched tracks of oblique incidence. These tracks are created by ions, whose charge and mass are more than those of nuclei of hydrogen. The composition of this component is mainly defined by ions with an energy of the order of 100 keV which are heavier than atomic ions of hydrogen. We call attention to the presence of the separate overetched oblique-incidence tracks which correspond, apparently, to heavy particles with an energy of the order of 100 keV/nucleon. The question about the composition of the slow track-forming component of the corpuscular emission of HD requires a separate investigation with the use of complicated analyzers, in which the detectors can be engaged only as an additional means.

For the more exact determination of the amount and the energy spectrum of the hydrogen component of the corpuscular emission of HD, we have counted the number of tracks on the detectors placed in the field of a

magnetic analyzer beyond the input diaphragm of $60\ \mu\text{m}$ in diameter remote from HD by 36 cm. In these experiments with the use of an external magnetic analyzer, the analyzed beam is ejected into the transanodic space along the diode axis. In this case, the possibility of a free dispersion of plasma in a specific direction is provided by the particular structure of the target. Such a configuration ensures also the possibility of the collimation of a separated beam at a large distance due to the weakening of the influence of the own magnetic field of a discharge on the motion of ions. In this case, the density of tracks produced by a beam is lower than the absolute one, and the direct count of tracks becomes possible.

In Fig. 6.3, we present a histogram with the results of counting the number of tracks on the sequence of intervals of the detectors corresponding to small changes in the cyclotron radius of ions $R = \sqrt{2E/Mm_p}/\varpi_B = 17\sqrt{E^*}$, where $\varpi_B = 8 \cdot 10^7 Z/M$ is the cyclotron frequency of ions, whose charge and mass numbers are, respectively, Z and M , in the field of a constant magnet of the analyzer with an induction of 0.8 T. The reduced energy of ions, $E^* = EM/Z^2$, is expressed in units of MeV. The histogram is constructed in such a way that the product of the ordinate and ΔE^* in the scope of each segment is equal to the corresponding number of tracks ΔN on the detector. Upon such a counting, all tracks are identified with the tracks of protons, for which $M = Z = 1$ and $E = E^*$. The counting of tracks for small cyclotron radii corresponding to the condition $E^* < 100\ \text{keV}$ was not performed.

Figure 6.3 illustrates a random energy spectrum of ions having passed through a small diaphragm under the indicated conditions after a single successful shot in units of E^* — “by recounting for protons” in the above-

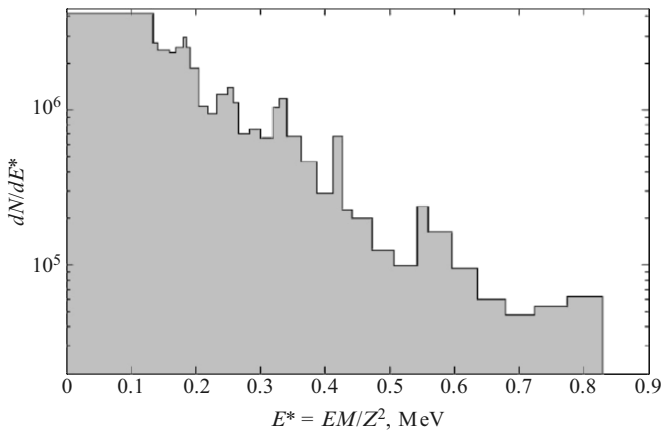


Fig. 6.3. Histogram of the spectrum of of a thin beam separated in the direction along the diode axis.

mentioned sense. The given spectrum includes $4.2 \cdot 10^5$ particles in all. Then, with regard to the geometric parameters of the experiment and taking assumption on the dispersion isotropy, we get the estimate of the order of $2 \cdot 10^{14}$ ions emitted by HD along all the directions. The mean value of E^* is equal to 0.27 MeV, and the total energy of the analyzed beam of ions in the region $E^* > 0.1$ MeV is $1.1 \cdot 10^5$ MeV. By recounting for “ 4π ” and assuming the dispersion isotropy, we get the integral contribution of HD to the emission of track-forming ions to be several joules.

In these measurements, we did not perform the analysis of the composition of a separated beam of ions over the whole spectrum (though such an analysis is possible). But, by the etching rate, the overwhelming majority of tracks is classified as those of nuclei of hydrogen and as those of protons on the high-energy edge. The presence of deuterons is registered in measurements by the same scheme with an extremely small diaphragm $\varnothing 5 \mu\text{m}$ in the main body of the distribution in the region of medium energies of the order of several hundred keV, where deuterons are close to protons by their number in the studied samplings (these results are discussed in Sec. 6.4.). In the low-energy region $E^* < 100$ keV, where tracks become extremely weakly pronounced and the packing becomes absolute, we expect the presence of heavier ions.

The given spectrum is of interest by its several features. It is characterized by the sharp edge in the region of high energies (in this case, about 800 keV for protons). The spectrum decays monotonically by a power law not inherent in the beam-related distributions. In this case, the great number of emitted protons and the presence of deuterons are strong arguments in favor of their nuclear origin.

6.3. Procedure of Track Analysis

The identification of particles, whose tracks are observed upon the etching of detectors, consists in the measurement of any two parameters of tracks sufficient for the determination of the charge and mass of the corresponding nuclei. The energy of particles incident on a detector is always of great interest. However, if the sort of particles is known, we can determine its energy, because the total path of particles belongs always to the number of the measurable parameters of a track. From the physical viewpoint, these three parameters define completely the process of ionization responsible for the formation of a latent track along the trajectory of a particle in the detector body. According to the modern ideas, the ionization-related properties of a particle of the given sort are phenomenologically characterized by the similarity of the distribution of the etching rate along a track and the dependence of the local rate of restricted energy losses (REL) of

a particle due to ionization on the residual path. Being expressed as the functions of residual path, both dependences acquire the universal character for ions of each type. Though the available theoretical models based on the notion of REL are, apparently, insufficient (this is shown by direct measurements of the local etching rate in Ref. 198 for ions heavier than α -particles), it follows from this analysis that a particle can be identified if two integral parameters are determined: for example, the path of a particle R (the length of a latent track in the detector body) and the thickness of the layer of a detector removed for the full time interval of the etching of a track

$$H = V_b \int_0^R \frac{dx}{V_t}, \quad (6.1)$$

where V_b and $V_t(x)$ are the rates of volume and track etchings, respectively.

The use of integral parameters is convenient for the measurement of tracks with paths of the order of $10\ \mu\text{m}$ which are observed as well-measurable objects after the 3 to 5-hour etching, being already overetched. For hydrogen tracks, such an approach is uniquely possible, because the nuclei of hydrogen do not create, in principle, deep and slanting tracks, and the determination of local values of the etching rate upon microscopic measurements in reflected or transmitted light seems to be impossible. The commonly accepted integral characteristic is the mean dimensionless relative etching rate of a track, V , namely, the ratio of the path length to the thickness of a removed layer:

$$V = R/H. \quad (6.2)$$

By determining the values of R and V in some way, we can determine, in principle, a sort of particles, because the parameters R and V are uniquely connected for particles of each sort, and the lines $V(R)$ for different sorts of particles do not cross on the plane (V, R) by joining in weakly split isotope families. The problem of identification of tracks is complex: it is necessary to possess the calibration data for a specific type of a detector and to ensure a sufficient accuracy of measurements upon the determination of values V and R for each individual track. The method of “loci” consists in the determination of that which groups of analyzed tracks belong to sufficiently narrow stripes in the plane (V, R) which are set the value and tendency of $V(R)$ to be amenable to the interpretation.

For the analysis of tracks with paths of less than and of the order of $10\ \mu\text{m}$, the TDS method (see Ref. 193) is the most convenient, because it is basically adapted to the processing of overetched tracks. The TDS method can be formulated on the basis of a simple analysis of the last phase of the

etching beginning from the time $t = t_R$, when the latent track is already etched, and only the volume etching of the detector substance is running. At the end of a conic track, a spherical rounding arises, whose radius grows by a simple law: $r = V_b(t - t_R)$ (the volume etching rate can be considered to be constant). This spherical surface is combined with the inner lateral surface of the track, which is formed by joining, according to the Huygens principle, the etching spheres originating from points of the surface. Upon further etching, the lateral conic component of a track surface decreases in height, and this surface disappears at the time $t = t_R(1 + V)$.

For this time interval, the surface layer of the detector h_R approximately equal to the path of a particle is etched, and the track loses all the signs of a conic pit and acquires the form of an ideal spherical segment. The overwhelming majority of the observed tracks is referred to such categories of "etching pits" which are the segments of ideal spheres. In this case, all the individual features of tracks are conserved only in the form of two parameters of spherical pits: the diameter and depth. But, in principle, this is sufficient for the determination of R and V . A remarkable feature of etching pits is the constancy of their depth L for the thicknesses of an etched layer $h > h_R$: it is clear that the lower point on the bottom of a spherical segment descends with the same velocity V_m upon the etching, as the plane surface of a detector does. It follows from simple geometric considerations that $L = R - H$ under the normal incidence, where H is defined by Eq. 6.1. This yields that the reduced etching rate for a track can be simply given in terms of L and R :

$$V = R/(R - L). \quad (6.3)$$

The increase in the diameters of spherical segments during the deep etching also obeys a simple law. The radius of the input hole of a track of the type of a spherical segment is some projection of the radius r of a sphere which is developing on the track bottom under the condition $h > h_R$. The radius r grows linearly depending on the thickness of the etched layer: $r = h - H$. But, in view of the constancy of L and obvious geometric considerations, it is clear that the linear dependence on h is also valid for the square of the radius of a spherical etching pit $(D/2)^2 = r^2 - (r - L)^2 = -2L(R - L/2) + 2hL$. The TDS method (see Ref. 193) consists in the measurement of two parameters of the law of a linear growth of the square of the diameter of a separate track, $D^2(h) = k(h - a)$, upon the successive etching of a detector in the phase where $h > h_R$ and the calculation of L and R by the simple formulas:

$$L = k/8, \quad (6.4)$$

$$R = a + k/16. \quad (6.5)$$

This allows us to find the reduced etching rate by Eq. 6.3:

$$V = (a + k/16)/(a - k/16). \quad (6.6)$$

It is basically important that the linear law for the square of a diameter becomes true for a finite time of the etching and is strict, rather than asymptotic. The data presented in Ref. 193 for nuclei heavier than the nuclei of hydrogen confirm that the dependence $D^2(h)$ really becomes linear shortly after that the thickness of the etched layer has exceeded the path of a particle. Indeed, the law of constancy for the depth of a spherical etching pit is satisfied sufficiently exactly. But the exactness of this linear dependence for a long-term etching of hydrogen tracks was not verified in Ref. 193. The advantages of the TDS method are its simplicity and the accessibility of the diameters of tracks for very exact optical measurements. As a drawback, we indicate the necessity of a long-term etching and multiple hourly successive measurements aimed at the verification of the attainment of the mode of linear dependence $D^2(h)$ and the enhancement of the exactness of determination of the parameters a and k . Just the last circumstance makes the TDS method to be labor-consuming and hampering the derivation of a rapid information in a series of current experiments.

But the main drawback of the TDS method is revealed upon study of hydrogen tracks, which seems to be, as was mentioned above, solely possible, because the investigation of tracks on the breaks of detectors taken from specific experiments is ruled out. All our measurements of hydrogen tracks show that the values of V determined by the TDS method are underestimated by 10% as compared to the theoretical ones calibrated on the data (see Refs. 194, 195) derived by the method of breaks. The reason for a decrease in the measured value V_{ex} consists in the systematic error of optical measurements of the diameters of tracks which is related to the underestimation of the diameter D_{ex} as compared to the real value of D by the value D_c introduced by the optical effect of displacement of the contrast boundary inside a track upon the observation of the track on a microscope in reflected or transmitted light (see Refs. 195, 196, 199). The value D_c is defined as a quantity which should be added to $D_{ex}(h)$ in the limit of a small thickness of the etched layer, $h \rightarrow 0$ in order to satisfy the condition $D_{ex} = 0$ (see Refs. 196, 199) and depends on the sort of a particle, its energy, and, possibly, on h for a large thickness of the etched layer. For protons of the MeV-energy range, $D_c \approx 1 \mu\text{m}$ (see Ref. 199).

The essential peculiarity of the TDS method is the following. Upon the determination of parameters of the limiting linear dependence of the square of the track diameter, $D^2(h) = k(h - a)$, a decrease in the measured diameter is proportional to the very value of $D(h)$. Obviously, this effect is

especially significant in the case of hydrogen tracks which are particularly small etching pits, for which we need to take into account the additional increase in the inner displacement $D_c(h)$ of the visible boundary of the very gently sloping edge of tracks. By determining the parameters a and k in Eqs. 6.5, 6.6 by two successive measurements of the track diameter, whose exact values have form of the sum $D = D_{ex} + D_c$, it is easy to get the following formula for the appearing error in the linear approximation in D_c :

$$a = a_{ex} - D_c D_{ex} / k_{ex}, \quad (6.7)$$

$$k = k_{ex}(1 + D_c / D_{ex}). \quad (6.8)$$

Thus, the correction for a needed to be made for the compensation of the effects related to D_c grows with the observed diameter of a track D_{ex} . In this case, a similar correction for k decreases, and it can be omitted. Eqs. 6.3–6.8 yield the expressions for the deviations of the ideal values of R , V from the measured ones:

$$V = V_{ex}(1 + D_c D_{ex} V_{ex} / 8R_{ex}^2), \quad (6.9)$$

$$R = R_{ex}(1 - D_c D_{ex} V_{ex} / 8R_{ex}^2(V_{ex} - 1)). \quad (6.10)$$

According to the signs on the right-hand sides of Eq. 6.10, it turns out that the correction of the experimental locus $V_{ex}(R_{ex})$ should be carried out by the passage to the lower values of R and the greater values of V . If we take the value $D_c = 1 \mu\text{m}$ and the diameter of an overetched track $D_{ex} = 14 \mu\text{m}$, then the relative measurement errors of V and R reach, respectively, 10% and 20% in the region of paths $R = 5 \mu\text{m}$, where the etching rate is close to the underestimated maximum $V_{ex} = 1.5$. The account of the additional increase in the value D_c in dependence on the etched layer thickness makes the corrections considerable else in the earlier measurements of diameters, by which we usually restrict ourselves. For the proper consideration of all the effects of D_c , it is insufficient to use only the linear approximation, and a further development of theory is needed with the purpose to improve the TDS method.

The other source of errors upon the determination of R , V within the TDS method is the randomness of the angle of incidence of a particle on the detector. The analysis shows (Ref. 193) that, at an arbitrary angle of incidence θ , the right-hand sides of Eqs. 6.5 and 6.6 determined from measurements define, respectively, $R \cos \theta$ and $V \cos \theta$. A more strict consideration shows that such simple answer is derived only for angles of incidence sufficiently close to 0° . Upon the analysis of the beams separated by diaphragms far from HD, all the tracks of the analyzed groups in the scope of submillimeter aggregates on a wide plane detector are characterized by a

common value of θ . Therefore, a deviation from the normal incidence does not affect the very form of a locus, and, to correct results of Eqs. 6.5 and 6.6, it is sufficient only to take into account the common coefficient $1/\cos\theta$ at small angles of incidence. In all the cases under consideration, this effect plays no significant role, and we can easily ensure the conditions, for which the account the angle of incidence gives corrections of about 1%.

Upon the measurements by the “rule of squares”, most exactly is determined the limiting depth of etching pits L . Our experience shows that the main source of errors is the inaccuracy in the determination of the thickness of an etched layer in each measurement cycle, which is related to the necessity to measure large lengths accessible for optical observations only at a magnification of 192 with an absolute error of $\pm 0.31\ \mu\text{m}$, whereas the measurements of the diameters of tracks are performed at a magnification of 960 and an error of $\pm 0.062\ \mu\text{m}$. The quantity L is determined with relative theoretical error $\delta_L = \frac{1}{2}\delta_D + \frac{1}{4}\delta_h \approx 2\%$. Upon the determination of a and R , the errors are doubled: $\delta_R \approx \delta_a \approx 4\%$. We note that the calculation of $V = R/(R-L)$ is affected by the presence of a difference in the denominator which can be small: the values of a path and a limiting depth of etching pits can be of the same order (i.e., this occurs for large values of V). Since the absolute errors satisfy the condition $\Delta L \ll \Delta R$, the relative error is

$$\delta_V = \delta_R + (\Delta L + \Delta R)/(R - L) \approx \delta_R(1 + V). \quad (6.11)$$

Thus, for large track etching rates, the relative errors of V increase proportionally to V . So if $V > 10$ (this is true beginning from carbon in the region of paths about $10\ \mu\text{m}$), the error of the determination of the mean etching rate of a separate track can be $\geq 100\%$. However, this fact says nothing about the boundedness of the TDS method used in the decoding of the tracks of heavy particles. Obviously, the very approach to the construction of “loci” for heavy particles in the coordinates (R, V) is not the best one: it is more convenient to construct the diagrams (R, V) directly, as it was verified upon the measurements of V as a function of the residual path in the underetched tracks of heavy particles from cosmic rays (see Ref. 200). Errors in the determination of the thickness of the etched layer can be compensated by measurements in a series of successive etchings and by the calculation of the optimum straight line $D^2(h)$ for every track by the least-squares method. The TDS method gives the best results in all the cases upon the determination of the limiting depth of a track L up to $10\ \mu\text{m}$ (moreover, it is solely possible if the method of breaks of detectors is excluded). The resolution power for experimental loci with close L and R can be ensured by measurements of a sufficiently wide collection of single-type tracks and by the determination of the mean statistical tendency of $L(R)$.

For the verification of the TDS method and for the purposes of practical calibration, we constructed a locus (R, V) of the short-length tracks of α -particles from a neptunium source closed by an aluminum foil of 10 mm in thickness. These measurements fill a lacuna in the reliable literature data. We mention the unique work (see Ref. 196), where the method of breaks of detectors used in the measurements of L and R gave the inadmissible variance of points in the region of the maximum $V \approx 4$ upon the recalculation to determine V , so that even a tendency of $V(R)$ cannot be guessed. In Fig. 6.4A we present the results of our measurements of V for the paths $R < 13 \mu\text{m}$ along with the values of V calculated on the basis of the data

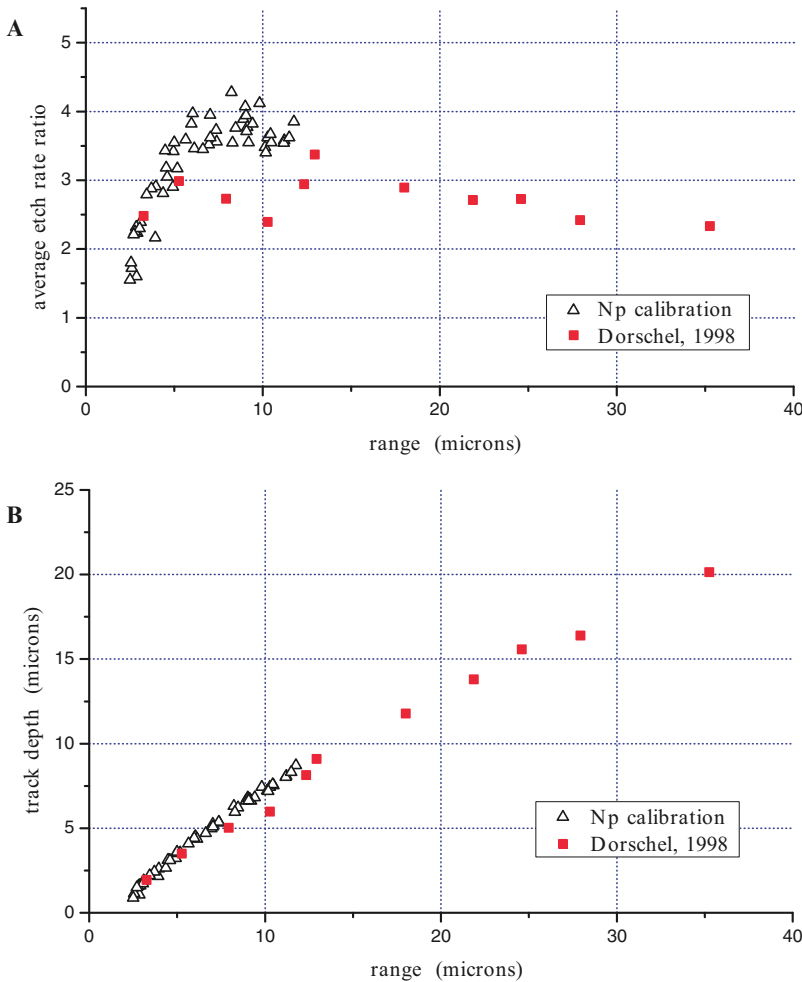


Fig. 6.4. Experimental loci of α -particles constructed by the TDS method on the basis of the data taken from Ref. 196: RV -locus (A); RL -locus (B).

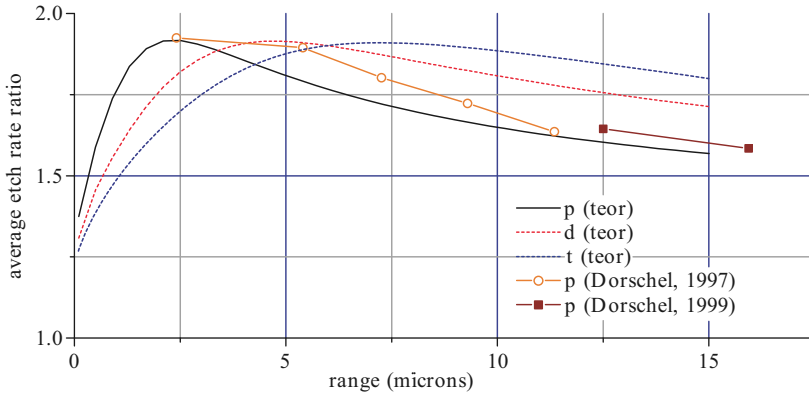


Fig. 6.5. Calibration of an RV -locus for isotopes of hydrogen by the data given in Refs. 194, 195.

given in Ref. 196. Here, each point corresponds to the measurements of 12 tracks, including those in the region of large paths. In Fig. 6.4B we present the locus (L, R) equivalent to Fig. 6.4A. Here, we also give the points corresponding to the data taken from Ref. 196 in the region of paths with maximum V . Figure 6.4 clearly demonstrates the advantages of the TDS method in the determination of $V(R)$ and $L(R)$ upon the study of short-length tracks.

To calibrate the loci of protons, we used the results given in Ref. 194, on the basis of which we succeeded to construct 5 points in the coordinates (R, V) in the energy interval from 0.2 to 0.7 MeV. In Fig. 6.5, we present the results taken from Refs. 194 and 195 along with the theoretical curves for all isotopes of hydrogen calculated within the model of restricted energy losses. These data do not reflect the effects of the critical diameter D_c , but remain a unique reliable reference point in the discussion of experimental data in the absence of reliable calibrations of the TDS method for nuclei of hydrogen.

6.4. Registration of the Image of HD on Track Detectors in an Ionic Obscure-Chamber and a Magnetic Analyzer

Ionic obscure-chambers, in which a track detector serves as a register of the image of a source of emission, are used in the studies of HDs of pulse discharges of the “laser” or “plasma focus” type (Refs. 201, 202). The scheme of an obscure-chamber used by us is shown in Fig. 6.6. The distance from the discharge axis to the input hole of the obscure-chamber is 15 mm, and the conditional coefficient of magnification of the obscure-chamber is usually equal to 2. Figure 6.7 gives the examples of the derived images.

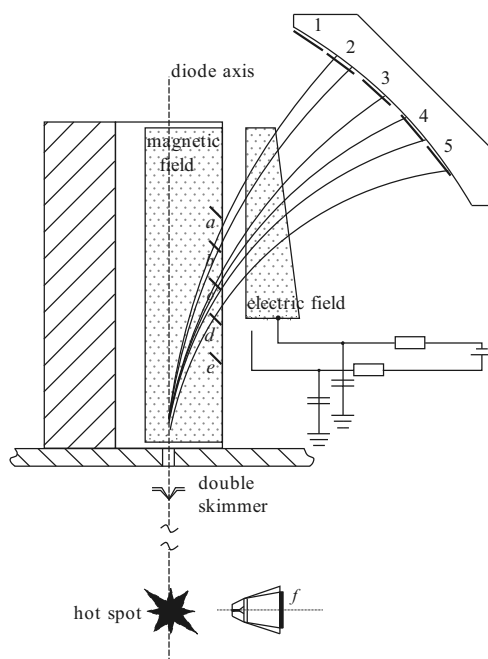


Fig. 6.6. Scheme of the registration of ions with track detectors. In the upper part, we display a Thomson mass-spectrometer, in which detectors CR-39 are in positions 1–5. A magnet with detectors *a–e* is used as a magnetic analyzer.

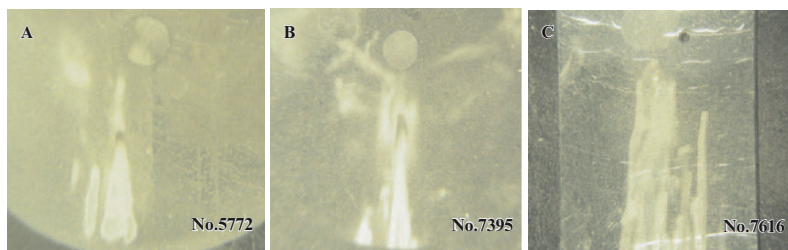


Fig. 6.7. Examples of the obscurograms of HD derived in ionic obscure-chambers. In the upper parts of images, we see a weakly pronounced round spot of 2 mm in diameter which is not a part of the obscurogram of HD and is formed by the plasma irradiation with the back side of the detector trough a tuning window in the back wall of the obscure-chamber. The spot center lies on the obscure-chamber axis.

The obscure-chamber is positioned in the equatorial plane of the diode passing through the conditional center of the diode or with a small angular deviation from the plane. The sighting of the obscure-chamber to a target is performed with the help of a laser beam directed along the

obscure-chamber axis through a hole on the back wall and a transparent detector CR-39. For obscure-chambers, we used diaphragms of 20 to 50 μm in diameter produced by the burning through of a tungsten foil of 20 μm in thickness with a laser. Less firm diaphragms are not suitable even for a single usage in our experiments.

The images of HD of the “laser” or “plasma focus” type (see Refs. 201, 202) are always observed in the form of a compact spherical source with a radius of several tens of microns. On the contrary, the track detectors in our obscure-chambers demonstrate a great variety of composite, as a rule, images of HD appeared after one shot. These images are represented by several stripes (and rarely by spots) formed by the dense aggregates of deep tracks which are vertically expanded (along the discharge axis). Moreover, each from these objects is characterized by single-type tracks.

In the scope of each stripe, the deepest and denser tracks are observed at the stripe beginning remote by several millimeters from the obscure-chamber axis. In this case, the image of the “head” of a stripe often looks to be enlightened in the transmitted light in a microscope due to a significant excess of the density of tracks over the absolute value, at which the overlapping of tracks begins. Stripes run under the detector edge remote by 10 mm from the obscure-chamber axis. In this direction (to the cathode, as compared to the diode axis), stripes spread strongly sometimes and shift to the horizontal (azimuthal) directions. In this case, tracks degrade in the scope of each stripe. Sometimes in especially contrast images on the periphery of “heads” of separate stripes, we observed scanty (about hundreds) local aggregates of especially long-length tracks. In such groups, we registered nuclear hydrogen tracks up to 40 μm in length.

The widths of separate stripes on track images are usually equal to 100 to 200 and more μm . The effect of horizontal splitting of stripes is seen on the images derived with the use of large 100 to 150- μm diaphragms. The analysis of the widths of separate stripes on images allows us to conclude that the diameter of HD is at most 100 μm , which is from $\frac{1}{5}$ to $\frac{1}{15}$ of the diameter of the used targets.

Obscure-chambers do not show any details of the region of corpuscular emission on the 10- μm scale. However, the image character testifies to the complicated process of emission from HD in the time limits of its luminosity. In essence, we see a single image in the form of a stripe or a cone which is represented by several fluxes of the fastest ionizing particles. These fluxes undergo splitting by the magnetic field of the diode which is present during the time interval of the dispersion of the fastest front of a plasma emitted by HD. The different times of motion of ions forming the closely located stripes in Fig. 6.7 from the start from the target cumulation

region to the detector and the difference of the angles of deviation from the obscure-chamber axis at the time of passing the diaphragm are explained by certain differences of the individual properties of ions (energy, mass, and charge) and by an instability of the magnetic field of the diode.

The time of flight of protons with energies of several hundreds of keV is of the order of 1 ns. In this case, the amplitude values of the magnetic field in the flight space between the discharge axis and the detector in an obscure-chamber, being of the order of 10 kGs, are sufficient in order that the angle of incidence on a plane detector on the obscure-chamber bottom be several degrees, which is observed experimentally. The vertical formation of stripes is explained by the broadening of the beams of single-type particles along cyclotron radii in the azimuthal magnetic field of the discharge current. Obviously, the velocity of particles of a beam varies in the limits of its emission time by other law relative to that of the field dynamics. In Fig. 6.7, tracks in the scope of each stripe degrade along the direction from the "head" downward, which corresponds to lesser cyclotron radii of ions on the periphery of detectors.

The horizontal splitting of stripes can be naturally explained by the action of the azimuthal electric field, which is parallel to the basic magnetic one, on the way of ions from HD to a detector on the obscure-chamber bottom. We observe the same effect allowing us to derive the "Thomson parabolas" in the analyzer of the thin beams of fast ions which uses the successively positioned zones of the transverse constant fields: magnetic and electric ones. In the real "self-organizing" experimental situation, these zones are superposed, and the fields are nonstationary. In this case, the nonstationary longitudinal component of the magnetic field of the diode comes into play, whose dynamics generates a vortex azimuthal electric field beyond the basic current channel of the diode, which leads to the horizontal splitting of stripes of the order of 0.1° in azimuth.

The most simple model of the formation of images in Fig. 6.7 is reduced, thus, to that HD emits the bounded collection of beams of single-type ions at every time, whereas the nonstationary screw magnetic field of the diode current imitates the operation of a Thomson analyzer of the beams separated by a diaphragm. Without the time-of-flight analysis of fast ions, the interpretation of track images in ion obscure-chambers cannot be unambiguous. Thus, considerable corrections to this scenario are possible.

In the frame of the proposed interpretation, we have to recognize the possibility of the "flickering" of HD, i.e., the repeated emission of the beams of single-type ions with identical energies in the scope of the total duration of luminosity of HD of the order of 10 ns. This is testified by the presence of equivalent stripes in different groups at different angular distances

from the obscure-chamber axis in multistripe images like those in Fig. 6.7B. In the scope of a “triplet” like that in Fig. 6.7A, such an effect is not observed. The “triplet” can be considered as an elementary image of HD in ion obscure-chambers for one more reason: it is seen as an elementary forming group in the multistripe images.

The described features of the obscurograms in Fig. 6.7 (we have totally got more than 200 especially informative contrast images) characterize the observed hot dot as an unusual, puzzling object. Such a characteristic is strengthened by the results of analysis of the composition of the beams of particles forming tracks on the multistripe images. In Fig. 6.8 we present the photo of HD convenient for analysis. (The horizontal lines in Fig. 6.8 are drawn for the convenience of separating the groups of tracks upon the implementation of microscopic measurements and are not a part of the track image). The practical complexity of analysis of individual tracks forming the separate aggregates consists in the necessity to measure the groups with the density of tracks which is less than the absolute one and is such that tracks do not overlap in the process of long-term etching. As usual, the stripes on images are characterized by the density of else underetched tracks which is higher than the absolute one. The choice of tracks on the lateral periphery of stripes is unsatisfactory due to both the overlapping of stripes on the image and the influence of the effects of angular scattering on the edges of a diaphragm. Such effects favor the expansion of stripes upon the construction of loci. In the case of the “triplet” represented in Fig. 6.8 stripes are well

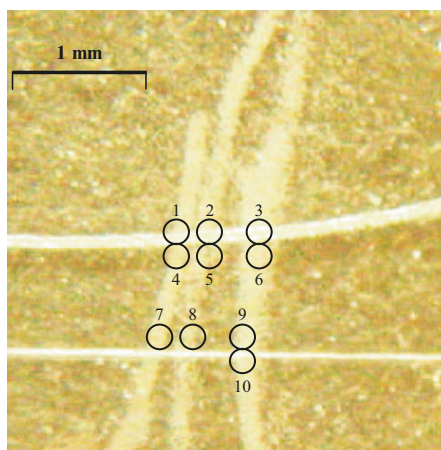


Fig. 6.8. “Triplet” as a typical image of HD derived on a track detector in the ion obscure-chamber with a $20\mu\text{m}$ diaphragm and the coefficient of magnification equal to 2 (horizontal lines are artificially drawn for the convenience of measurements and are not a part of the image of tracks; circles mark the area, where the groups of tracks were separated for measurements).

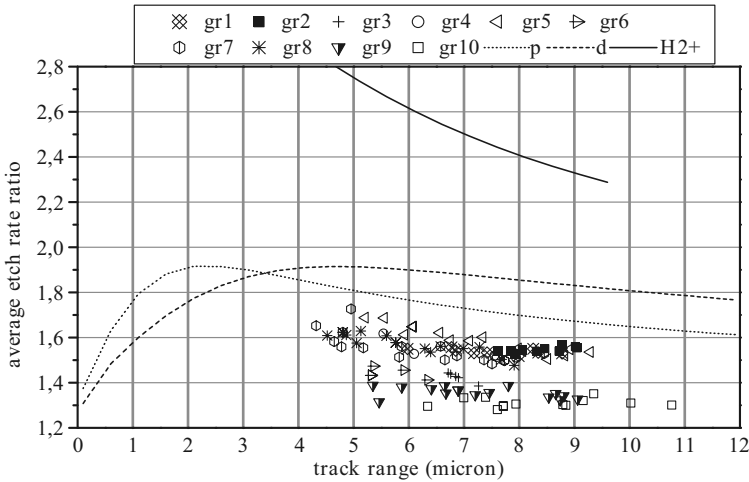


Fig. 6.9. A locus of 100 tracks taken from 3 stripes of “triplet” No. 6512. Here, we present also the theoretical curves for protons (a black continuous curve), deuterons (a dashed one) and molecular ions H_2^+ (the upper eroded curve in the range of paths from 5 to $10\mu m$).

separated and are characterized by a density of tracks which is convenient for the measurement of diameters by the TDS method.

In Fig. 6.9 we show an RV -locus of 100 tracks from ten groups separated in the limits of three horizontal lines on the image given in Fig. 6.8: groups 1, 4, and 7 are in the left stripe, groups 2, 5, and 8 are in the middle one, and groups 3, 6, 9, and 10 are in the right one. A remarkable feature of this locus is the presence of the splitting into two stripes. It turns out that the lower stripe on the locus is formed only by the tracks from groups 3, 6, 9, and 10 from the right stripe in Fig. 6.8, whereas the remaining tracks of the groups taken from the left and central stripes in Fig. 6.8 form the upper stripe of the locus. Because all the tracks were processed and measured according to the TDS method in the same manner, and we cannot indicate any physical reasons for the splitting of a locus of identical particles with identical paths, it is natural to assume that the tracks presented in Fig. 6.8 belong to two different sorts of particles. In this case, we may choose only among the nuclei of hydrogen, because only they are characterized by the etching rate satisfying the condition $V < 2$ in the region of paths from 4 to $12\mu m$ represented on the last locus.

The comparison of the theoretical curves $V(R)$ for the isotopes of hydrogen (see Fig. 6.5) with the experimental locus shows that the measured values for the lower stripe in Fig. 6.9 turn out by 10% less than the lowest known values of V set by the theoretical calibrated locus for protons in Fig. 6.5 (e.g., $V = 1.8$ for a $10\mu m$ path of protons, whereas the lower stripe

in Fig. 6.9 approaches 1.65). The splitting of stripes on the experimental locus reaches 10% and corresponds quite closely to the theoretical separation of curves $V(R)$ for protons and deuterons. But, on the whole, the measured locus turns out to be shifted downward by the same 10% relative to $p-d$ lines in Fig. 6.5. A shift of the experimental values of $V_{ex}(R)$ determined by the TDS method downward relative to the calibrated hydrogen loci was observed by us upon the analysis of all track aggregates systematically, including those on the detectors in obscure-chambers.

In the previous section, we have discussed two possible reasons for that the values of V and R determined by the TDS method can be underestimated. One reason consists in the following. Upon the oblique incidence at small angles of incidence θ , experimental values are the projections of the total quantities: $V_{ex} = V \cos \theta$, $R_{ex} = R \cos \theta$. In the case under consideration, the influence of the angle of incidence should be excluded: all measurements on the detector shown in Fig. 6.8 are performed on a submillimeter area remote by 4 mm from the obscure-chamber axis. For the flight base of 30 mm between the diaphragm and the detector, a deviation of the trajectories of ions upon the passage of the diaphragm gives $\cos \theta = 0.99$, because the motion inside the obscure-chamber is rectilinear. The last circumstance was checked by us experimentally. Thus, we cannot explain the 10% shifts and the splitting of stripes on experimental loci by the influence of the angle of incidence.

The second reason lies in that the error of the determination of R and V increases during a sufficiently long etching upon the measurements of the track diameter visible in a microscope as compared to its real value. In Sec. 6.3., we have discussed the consequences of this effect for the algorithm of the TDS method, especially upon the etching of hydrogen tracks, for which we have to take into account the increase in the internal shift $D_c(h)$ of the visible boundary of tracks, being the relatively small etching pits. The values of corrections can reach 10% and 20%, respectively, for V_{ex} and R_{ex} .

Without the strictly determined calibration corrections for the loci of protons and deuterons constructed according to the TDS method, we restrict ourselves by the indication of the necessity of a congruent displacement of the experimental locus $V_{ex}(R)$ upward to the coincidence with similar theoretical curves $V(R)$. At the coefficient of similarity ≈ 1.1 to 1.2, we observe a good coincidence of a pair of track stripes in Fig. 6.9 with the theoretical curves for protons and deuterons with paths from 4 to 12 μm . (It is worth noting some increase in the splitting of track stripes as compared to the theoretical value, which is obviously related to the increase in D_c for deuterons as compared to D_c for protons according to the general tendency (see Ref. 199). This circumstance should be referred to the advantages of the TDS method).

With the purpose of the most exact measurement of tracks in the beams with a deliberately weakened density, we undertook the construction of a locus of tracks on the detector placed beyond a limitedly small diaphragm of $5\mu\text{m}$ in diameter on the diode axis beyond the anode in the magnetic field of an external magnetic analyzer. The passage of a beam along this direction was provided with a special construction of a target. In this case, the influence of the own magnetic field of the diode current was minimum. The diaphragm was placed at the input to the magnetic analyzer at the distance of 36 cm from HD, and the track detector was at a distance of 30 mm beyond the diaphragm in a constant magnetic field of 0.8 T. In this case, the detector plate was bent to provide the normality of the incidence of ions moving along the cyclotron orbits in the magnetic field of the analyzer.

In Fig. 6.10A we show the pattern of the distribution of tracks on the detector in experiment No. 7496. The density of tracks is extremely low, which creates the best conditions for measurements. But even in this case of

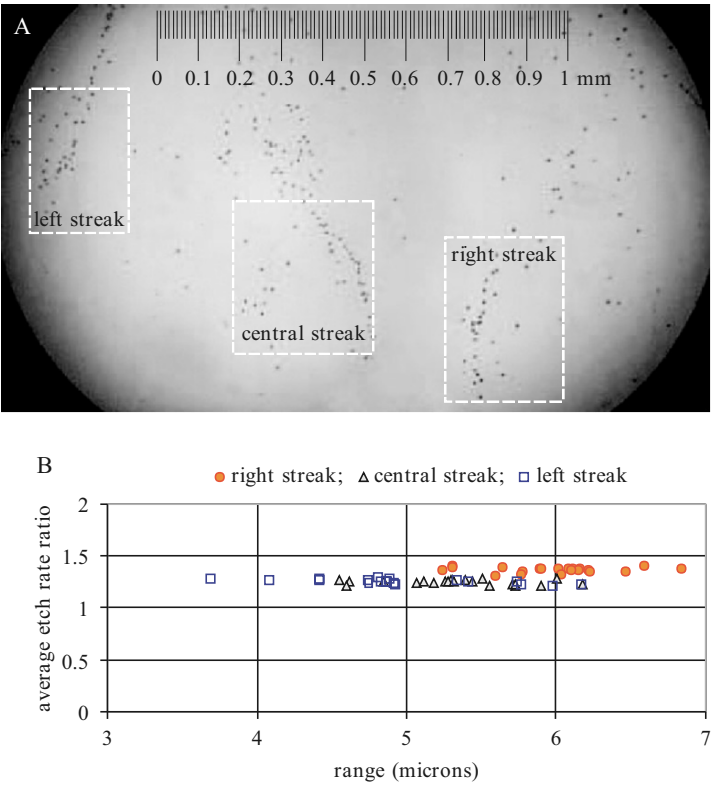


Fig. 6.10. Aggregates of tracks on a detector beyond the $5\text{-}\mu\text{m}$ diaphragm in a magnetic analyzer (A); locus of the separated groups of tracks (B).

the extremely low packing, we see three basic tracks on the detector, which corresponds to the observations in the obscure-chamber and confirms the above-presented considerations on the role of a magnetic field being external relative to HD as for the splitting of the beams emitted by it. For the analysis, we separated 80 tracks located arbitrarily over the whole detector with a preferred attachment to the aggregates of tracks. The corresponding locus is shown in Fig. 6.10B. It demonstrates the presence of splitting like in the above-mentioned case of the registration on the detector in the obscure-chamber.

In the last case, tracks are characterized by lower paths and are more overetched for the 6 to 12-h cycle taken by us within the TDS method for hydrogen. Tracks in Fig. 6.10B can be identified as those of protons and deuterons, but with the following additional remarks. These stripes are very close one to another and are separated only due to both the attained 4% accuracy in the determination of the RV -coordinates of separate tracks and the shift of a lower stripe to the side of lesser paths. This shift testifies to that, namely, the lower stripe on the given locus corresponds to deuterons, because, with regard to the deviation of particles in the magnetic field of the analyzer, the discovery of deuterons and protons in the scope of the common 1-mm area is possible at lower energies and paths of deuterons as compared to those of protons. However, the ratio of energies $E_p : E_d = 1 : 2$, which must be valid for protons and deuterons registered jointly under the complete coincidence of cyclotron orbits, is not strict in the given case where the analyzer is used in the obscure-chamber geometry with one diaphragm at the input.

We mention two reasons defining the difficulty to carry out the energy-involved analysis in the case where a magnetic analyzer is used in the optical geometry of an obscure-chamber without a collimator for the derivation of split stripes of tracks similar to those observed in an obscure-chamber. The optical effect inherent in an obscure-chamber allows us to reveal the splitting of beams shown in Figs. 6.8 and 6.10A due to the angular deviation of beams upon the passage of the diaphragm which occurs under the influence of the electromagnetic field near HD. Since the cyclotron orbits in the external magnetic field inside the analyzer are very sensitive to the presence of transverse perturbations of the velocity of particle at the input of the analyzer, we cannot indicate a value of the cyclotron radius of ions registered at a given point on the detector with a sufficient accuracy. Figure 6.10A gives an example of the successful combination of the effects of the magnetic fields of the diode and analyzer. In this case, the analyzer is mainly used to decrease the density of tracks.

According to the discussion given in the previous section, we have also to correct the locus shown in Fig. 6.10B for the account of corrections related to the errors of optical measurements of the diameters of tracks. This pro-

cedure leads to an increase in V_{ex} and to a decrease in R_{ex} , and the relative decrease in R_{ex} is approximately by the factor $1/(V_{ex} - 1) \approx 2$ less than the increase in V_{ex} . It is possible than the splittings of loci in Fig. 6.10B reflects the lucky circumstance that a defect D_c of the experimental measurement of the diameters of tracks depends on the sort of particles (see Ref. 196) and is relatively large for weakly pronounced deuteron-induced tracks. After the correction, the ratio of the energies of protons and deuterons on the proper locus approaches 1 : 2. But the vertical separation of the corrected loci becomes insignificant, because the theoretical curves $V(R)$ for protons and deuterons practically coincide in the region of paths $R \approx 4$ to $6 \mu\text{m}$.

The most important result consists in that deuterons are also presented in Fig. 6.10A by two basic left stripes of tracks, and protons are presented by a single right stripe like in the above-considered case of a detector in the obscure-chamber in Fig. 6.8. In the region of paths under consideration corresponding to the energy range 0.35 to 0.5 MeV, the ratio of the numbers of protons and deuterons is also equal to 1 : 1 by the order of magnitude (it would be incorrect to give an exact relation for the given random sampling of tracks).

The indicated nuances are of importance upon the interpretation of experimental data and should be confirmed under the subsequent development of the TDS method and the comprehensive calibration of loci for nuclei of hydrogen within this method. Nevertheless, the above-presented consideration yields the following main conclusion: the employed method can reveal the fine effect of splitting of the loci of the isotopes of hydrogen. The derived data testify to the rightful presence of deuterons together with protons in the corpuscular emission from HD in our experiments.

We emphasize for the subsequent discussion that the above-presented data are derived in the experiments with the lead targets possessing the extremely low solubility of hydrogen, so that the observed amount of fast ions of hydrogen exceeds the rated limiting content of hydrogen in the damaged region of a target by several orders in magnitude. It is worth also noting that the residual gas in the chamber did not contain a hydrogen blend enriched with deuterium.

6.5. Measurements of Tracks with a Thomson Mass Spectrometer

In the connection to the above-presented interpretation, the question arises: Why do we exclude the possibility to consider tritium instead of deuterium as a component of a doublet on the loci constructed above, being additional to protons? Indeed, we may consider the appearance of the tracks of tritons in the form of stripes with the density close to that of the tracks of protons

in some region of parameters as improbable. But the calibration and the accuracy of the construction of loci are not so reliable in order to exclude such a possibility. This question is answered by the studies performed with the use of a Thomson mass-spectrometer, whose first results are presented in this section. It turns out that, indeed, the tritium component is present in the emission of HD in our experiments. But its amount is, apparently, by several orders less than the total amount of deuterons. Moreover, together with tritons, *we registered the tracks of one more type of unstable light nuclei, namely ${}^6\text{He}$* . A Thomson analyzer helps us to clarify some peculiarities of the energy spectra of ions of the hydrogen family and to confirm that the content of deuterons in the flux of fast protons is increased as compared to the natural one.

The scheme of a Thomson analyzer is shown in Fig. 6.6. Its structure is based on a magnet with the area of poles $5 \times 20 \text{ cm}^2$ and with a field of 8.0 kGs in an 8-mm gap used in other measurements as a magnetic analyzer. Figure 6.6 keeps the proportions of various units of the mass-spectrometer. The electrostatic field in the 8-mm gap between electrodes of the trapezoidal form reaches the value $E_{\text{max}} = 7.25 \text{ kV/cm}$. The analyzed beam is separated at the input of the magnet gap by means of a double skimmer with an inner diameter of 1 mm. In the electric field, ions also enter through a 1-mm split. The analyzer is evacuated autonomously up to 10^{-5} Torr.

As a peculiarity of the use of a Thomson analyzer in the given configuration, we mention the refusal from a two-diaphragm collimator present in the classical scheme (see, e.g., Ref. 203). In essence, our Thomson analyzer is a large obscure-chamber provided with internal static fields, where the maximum distances of detectors and the input diaphragm-skimmer from HD were, respectively, 66 and 36 cm. To ensure the output of a plasma in the transanodic region where the input diaphragm of the analyzer was positioned, we used the cylindrical targets located normally to the cathode – anode interval in experiments with a Thomson analyzer. In this case, we noted that if the target axis is directed normally to the magnet gap (in parallel to the magnetic field), then the detectors in the analyzer are abundantly filled by tracks. Upon the location of a target in parallel to the magnet gap, tracks on the same detectors are completely absent. In our opinion, this is explained by that the particles escaping from the target center move along trajectories slightly twisted by the electromagnetic field relatively the target axis. Therefore, entering the magnet gap, they “fall” onto the surfaces of the magnets.

The use of a strong magnet with a large area of poles makes our analyzer comparatively compact and suitable for the resolution of parabolas in the energy range of ions $E^* = (M/Z^2)E = 0.3$ to 2.3 MeV (we recall that

the parameter E^* is a characteristic of the deviation of ions with energy E in the magnetic field of the analyzer; in our case, the cyclotron radius of ions in the magnet gap is given by the relation $r = 18.0 \cdot \sqrt{E^*}$ cm; and E and E^* are taken in units of MeV). The deviation of ions in the electric field is unambiguously defined by r or E^* if the coefficient of proportionality M/Z is known: $= (M/Z)f(E^*)$. This relation takes the form of "Thomson parabolas", if we introduce the coordinate x marking the deviation of ions in the magnetic field on detectors instead of E^* (see Fig. 6.11a). A calibration of the characteristic f of the analyzer was performed by protons, for which $M/Z = 1$ and the deviation from the plane $= 0$ (the plane of the split) is minimum for every value of E^* or x . In this case, the locus of deuterons and all the ions, for which $M/Z = 2$, is defined by the strict doubling of the coordinate of a proton Thomson parabola against every value of x . Beginning from $M/Z = 2$, the parabolas can be filled by multicharge ions. Similar to mass-spectrometry, in order to completely identify an ion, we have to solve the problem of additional determination of its charge number, i.e., a value of Z for the given trajectory, and its atomic number defining the name of the corresponding chemical element. Such a basic possibility is given by the TDS method used by us. This method is especially efficient for the analysis of individual characteristics of the tracks of light ions.

In Fig. 6.11a we present the photo of parabolas registered after the first successful shot No. 8527 derived with the use of a Thomson analyzer on five detectors of $2.5 \times 2.5 \text{ cm}^2$ in area etched 4 h under standard conditions. The parabolas are seen by a naked eye on the sections where the density of tracks is sufficiently high and the craters of tracks are united. In a microscope, we see also other segments of parabolas and other parabolas in the places where the density of tracks is less than the absolute one. On the whole, we see about 20 parabolas. Even in the given structure of the mass-spectrometer intended for the registration of ions only in the approximate range $0.3 < E^* < 2.3 \text{ MeV}$ where only a part of the magnet perimeter is used, it is a high value as compared to other analyzers known from the literature.

The $M/Z = 1$ parabola formed by protons is easily recognized as most close to the zero line $y = 0$ and most intense. The corresponding stripe begins on detector 4 against the coordinate $x = 9.5 \text{ cm}$ corresponding to the energy of protons $E = E^* = 0.8$. We note that a small group of proton-induced tracks is also observed on detector 3. This calibration is derived by using the rapid procedure of calculations of R by the thickness of an etched layer H , at which the etching of the latent track is terminated. This time moment is easily determined by the beginning of the enlightening of craters during the repeated cycles of etching. According to the basic formulas Eqs. 6.1, 6.2 in

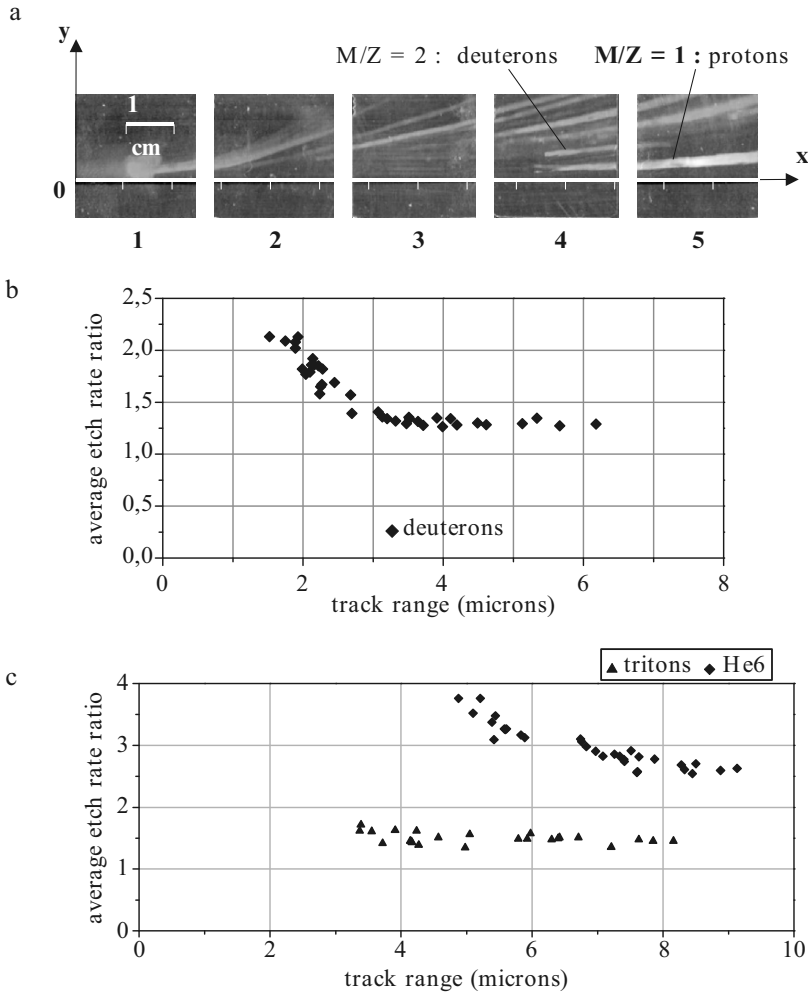


Fig. 6.11. Thomson parabolas on 5 detectors of $2.5 \times 2.5 \text{ cm}^2$ in size (experiment No. 8527) (a); locus of 40 tracks of the “head” of the $M/Z = 2$ parabola (b); locus of 50 tracks of the “head” of the $M/Z = 3$ parabola (c).

Sec. 6.3., a path is defined by the equality $R = VH$. For protons with paths $> 10 \mu\text{m}$, we can use the approximate formula $R = 1.5H$ and find then the tabular value of the energy of a proton by a certain path. The procedure of correction of this result in subsequent iterations with the use of the calibration dependence $V(R)$ in Fig. 6.5 gives a 20% correction for the rated value of E^* derived without regard for the edge effects of the magnetic field in the analyzer. The exact calibration of the offset $x(E^*)$ by protons in the whole range $x = 3$ to 13.5 mm for our analyzer is still to be performed.¹ We did

¹On the first stage of studies, the obstacle for solving this problem is the threat to lose

not pose a purpose to determine the energy spectra of particles forming the separate parabolas in Fig. 6.11a in order to get separate components of the spectrum shown in Fig. 6.3 in Sec. 6.2. which is integral by the sorts of particles. The result of the local calibration of $x(E^*)$ should be understood in the following way: there are no doubts in that we registered mainly protons in the coordinates of the $M/Z = 1$ parabola indicated in Fig. 6.11a.

We consider the identification of protons so comprehensively because there is no difficulties with the identification of deuterons in Fig. 6.11a. The $M/Z = 2$ parabola is unambiguously determined by a simple doubling of the ordinates (x) of the $M/Z = 1$ parabola is represented in the figure by the tracks with a packing close to the absolute one in the scope of corrected values $E^* = 0.45$ to 0.9 MeV on detectors 4 and 5. This parabola is present also on detector 3 up to the value $E^* = 1$ MeV with a loose packing of tracks. Just this segment is convenient for the analysis of tracks by the TDS method which is simplified in this case, because we must not take care of a high accuracy of the construction of a locus. Indeed, a thin splitting of a locus of the tracks taken from the same parabola is impossible.

In Fig. 6.11b we show the result of the rapid construction of a locus of the group including 40 tracks (the $M/Z = 2$ parabola) by three consequent measurements of the diameters of overetched tracks. This locus indicates unambiguously that the tracks belong to the hydrogen family, since the values of the mean etching rate are in the interval $1 < V \leq 2$ (see Fig. 6.5). A single isotope of hydrogen satisfying the additional condition $M/Z = 2$ is a deuteron. Those α -particles which are characterized by the same value $M/Z = 2$, but belong to the nearest helium family of light nuclei, demonstrate the ionization rate V in the region of paths indicated in Fig. 6.11b which exceeds twice the values of V for the nuclei of hydrogen (see Fig. 6.4a). We note also that the paths from 2 to $6 \mu\text{m}$ presented in Fig. 6.11b are characteristic of deuterons with energies of 0.2 to 0.45 MeV. This interval corresponds to the reduced energy $E^* = (M/Z^2)E = 0.4$ to 1.0 MeV, because $M = 2$, $Z = 1$ for deuterons. Just this interval corresponds to the positions of tracks chosen for the construction of a locus in Fig. 6.11b.

The $M/Z = 2$ parabola contains also single separated large tracks, the rapid analysis of which showed that they are characterized by the ionization rate $V > 10$. They are tracks of heavier, completely ionized ions (with mass not lesser than that of carbon) with energies of several hundreds of keV per nucleon. These tracks do not form an explicit stripe, are in the space between parabolas, and can be produced by scattered fast heavy ions, whose signal is distorted by the analyzer because of some reason.

a valuable information given by the $M/Z > 1$ parabolas containing short-path tracks as a result of the application of a many-hour etching long-path proton tracks according to the TDS method.

The $M/Z = 3$ parabola is well represented on detectors 4 and 5. In this stripe, we saw the aggregation of tracks of several types, whose analysis is hampered, however, due to a large density of tracks. Apparently, the tracks with $V > 10$ dominate in this case. This parabola is well-seen in a microscope and on detector 3, where it is exactly determined by the position of the “head” of a parabola of protons, i.e., by tripling the central corresponding coordinate. However, the $M/Z = 3$ parabola almost reaches the left edge of detector 3, where both the proton and deuteron parabolas are empty. On this section, we observe the tracks corresponding to two sorts of light ions with small density. It is obvious prior to the measurements on a microscope.

The results of the construction of a locus for the track sampling from the parabola “head” to E^* equal approximately to 1 MeV are presented in Fig. 6.11c. The lower stripe on the locus testifies to that one of the sorts of tracks of the $M/Z = 3$ parabola belongs to nuclei of the hydrogen family: the locus satisfies the criterion $1 < V < 2$ and very close to the hydrogen $V(R)$ -dependences (see Fig. 6.5) and to the locus $M/Z = 2$ presented in Fig. 6.11b. Like the previous case of the identification of deuterons, we must not unite the experimental loci on one diagram in order to see the effect of isotope splitting. For the final identification of tracks in the scope of the parabola, its index M/Z serves a more reliable criterion. In this case, it is sufficient to restrict oneself by the 10% accuracy of measurements realized by us in the rapid TDS procedure upon the analysis of detectors exposed in a Thomson analyzer². Like the case of the identification of deuterons, we can unambiguously conclude that the hydrogen tracks in the scope of the $M/Z = 3$ parabola, whose locus is presented by the lower stripe in Fig. 6.11c belong to the heaviest isotope of hydrogen, to tritium.

The potentialities of a Thomson analyzer are not exhausted by the identification of protons, deuterons, and tritons in the scope of the first three parabolas in Fig. 6.11a. The comparison of the upper stripe in Fig. 6.11c with the experimental $V(R)$ -dependence for α -particles (see Fig. 6.4) indicates that the corresponding tracks belong to a isotope of helium. In the region of paths 5 to $10\ \mu\text{m}$, the basic criterion for the etching rate, $2 < V = 4$, is satisfied because the difference of $V(R)$ -dependences in the scope of each isotope family is reduced to some shift of maxima depending on paths. This

²Upon the construction of “loci”, we do not indicate, as always, the errors of the determination of R and V for individual points, because the result of calculations depends not only on the accuracy of the procedure of successive etchings and the application of the least-square method upon the construction of the asymptote according to the TDS method, but on the fluctuations of the physical parameters of tracks, to which the accuracy of measurements is not referred. The final error of the determination of the tendency of $V(R)$ is reflected in the statistical broadening of the stripes representing the loci of tracks in the coordinates R and V .

effect is shown in Fig. 6.5 for hydrogen. The single isotope of helium satisfying the condition $M/Z = 3$ is nucleus ${}^6\text{He}$.

We note that the range of paths of tritons indicated in Fig. 6.11c corresponds to the interval of their energies $E = 0.35$ to 0.7 MeV, which corresponds to the interval $1.1 < E^* < 2.1$ MeV, because $E^* = 3E$ in this case. For nuclei ${}^6\text{He}$, we have $E = (Z^2/M)E^* = (2/3)E^*$, and their tracks have lengths of about $8\mu\text{m}$ according to Fig. 6.11c. Hence, they must have energy of at most 1.4 MeV. With regard to the fact that α -particles with a path of $6\mu\text{m}$ have energy of 1.3 MeV, we convince ourselves that the upper locus in Fig. 6.11c represents, indeed, particles from the family of helium. Thus, even without exact calibrations of $E^*(x)$ for an analyzer in the high-energy region and without tabular values of the paths of nuclei ${}^6\text{He}$ in a CR-39 detector, we *have unambiguously identified nuclei of tritium and nuclei ${}^6\text{He}$ in our analyzer.*

The total number of tracks in the scope of the 3rd parabola on detector 3 (in the region under study) is estimated at the level of 1% of the number of deuteron tracks in the region of the absolute packing on detector 4; the last number equals, in turn, approximately 1% of the number of tracks of protons on the $M/Z = 1$ parabola.

To get the additional confirmation of the fact that the particles registered by detectors are created namely at HD of a target, we performed two testing experiments. In the first case, we used a copper disk of 15 mm in diameter and 0.2 mm in thickness as a target. The discharge was directed onto the disk center. Such a geometry of the target did not allow a collapse to develop in its volume, though the target was damaged. After the execution of this experiment, all detectors in the analyzer were completely empty. In the second experiment, the disk was replaced by a copper stripe of the same thickness. Its axis was in the position favorable for the derivation of tracks (normally to the magnet gap). The result of the experiment was the same: we found no tracks on the detectors.

6.6. Observation of Nuclear Tracks

In a number of experiments, we observed, in addition to a continuous set of plasma tracks, the enhanced number of nuclear tracks close or identical to α -particles by the track etching rate. This effect is qualitatively recognized as some increase in the natural α -background accumulated by detectors for the time intervals of the preparation and the execution of measurements for each shot. But, in some experiments, we observed a local sharp increase, by four to five orders above the background, in the number of α -tracks on the detectors positioned near HD (at a distance of 1 to 4 cm) and shadowed from the direct plasma irradiation.

Such a situation arose firstly in some experiments with an ion obscure-chamber. In this case, we determined by the incidence angle on a detector that a source (or sources) was in the majority of cases not near the detector surface (as upon the registration of the natural background), but at some distance from it. We also observed that chaotically directed α -tracks appeared most frequently upon the use of diaphragms with a maximum diameter of about $300\text{ }\mu\text{m}$. However, the highest number of α -tracks was observed in the experiment, in which the detector was positioned in the transanodic chamber at once beyond the holder of a target beyond a $\varnothing 25\text{-mm}$ screen with 30 holes in the form of 8 radial rows which let pass about 1% of the plasma flux from a damaged target. The pattern of the chaotic filling by tracks and their individual peculiarities seen under a greater magnification are indistinguishable from α -tracks derived as a result of the 30-s exposure of the detector above a Pu^{239} source with an intensity of $10^4\text{ pulse/cm}^2\cdot\text{min}$. In our case, however, the irradiation of the detector occurred for the plasma expansion time in the channels of the anodic screen (about 100 ns). This is testified by the local presence of track aggregates only in the scope of 2 from 6 sectors beyond the circular screen and their practical absence on the other detectors positioned in the transanodic chamber in the same experiments.

It is natural to assume that the reason for the appearance of a great number of α -tracks is the presence of short-life (about 10^{-7} s) α -radioactive products in the plasma of a damaged target. This explanation seems to be most acceptable for the interpretation of the chaotic track aggregates.

However, in some cases, the observation of a growth of the α -background is added by the registration of collective centralized families of tracks similar to the tracks of α -particles. We note that the former can be referred neither to the natural α -background, nor to an anomalous one. In Fig. 6.12, we give the example of a local 4-track “bush” including one vertical and three slanting tracks formed by particles, whose ionizing abilities are equivalent or close to those of α -particles. The image is derived upon the contrast means of observation of tracks by the focusing of a microscope on the optical axes of slanting tracks. Similar families were observed on detectors on the inner lateral surface of a transanodic chamber at a relatively large distance (about 2 cm) from the discharge axis. We note that the appearance of “bushes” formed by four α -tracks is impossible under the ordinary background α -activity of Rn even in the hypothetical variant of the fixation of a nucleus Rn^{222} and its progeny on the surface of a detector at some point.

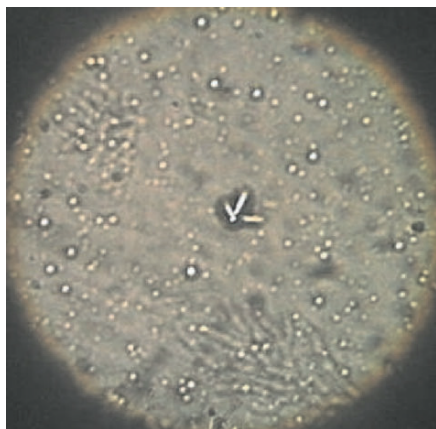


Fig. 6.12. Track cluster containing 4 tracks: central vertical and 3 slanting ones.

Moreover, our systematic observations of the natural α -background did not reveal even a three-track “bush”.

The observations of centralized track clusters are referred to rare events. All the more, the data of the registration of unique track clusters including the great number of tracks joined around one clearly separated center deserves a significant attention. The observations of centralized track clusters are referred to rare events. All the more, the data of the registration of unique track clusters including the great number of tracks joined around one clearly separated center deserves significant attention.

The example of such a family including 20 tracks is given in Fig. 6.13 (shot No. 3737). In this case, a detector was positioned in a collimator at the input to the magnetic analyzer for the control over a plasma beam distinguished on the front of a plasma expanding from a damaged target. Its distance from the discharge axis was about 12 cm. The family shown in Fig. 6.12 was registered at a distance of 1 mm from the edge of a 3-mm hole letting a beam pass. The hole edge and the signs of other nuclear clusters are seen in Fig. 6.14 on the same detector. The analysis of the track cluster shown in Fig. 6.13 indicates that it is formed by α -particles with energies up to 10 MeV and possibly by other light nuclei emitted from a single center at a distance of about $220\text{ }\mu\text{m}$ from the detector surface. The dispersion directivity diagram testifies to the fact of the instantaneous decay of a heavy nucleus, whose velocity was of the order of that typical of the motion of the particles formed the tracks.

A unique track cluster from 276 tracks was registered after shot No. 4109 on a detector positioned at once beyond a grid screening the collectors at the output of the magnetic analyzer (the experiment scheme is

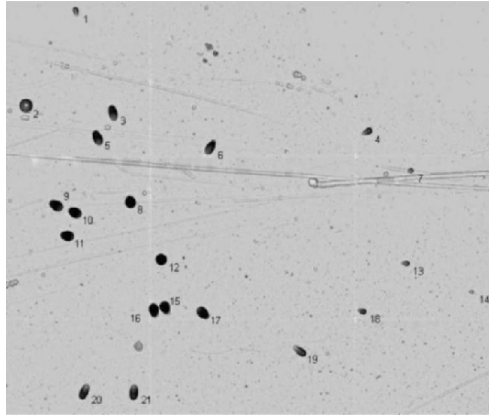


Fig. 6.13. Cluster from 20 tracks formed by the particles emitted by a moving center at a distance of $218 \pm 8 \mu\text{m}$ above the detector surface. The image is derived with the help of a scanning microscope Olympus. The diameter of the central track No. 12 is $23 \pm 3 \mu\text{m}$.

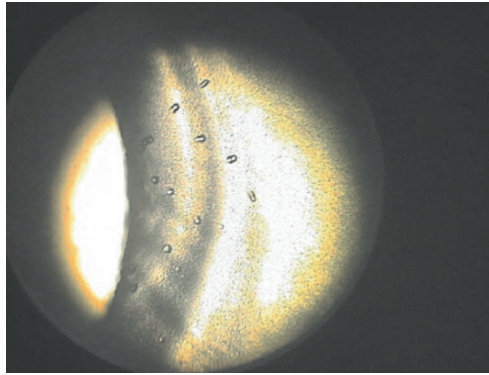


Fig. 6.14. Nuclear tracks on the edge of a hole letting pass the plasma beam into a magnetic analyzer with the same detector as in Fig. 6.13. Nuclear tracks are seen against the uniform background of small plasma tracks and are partly joined into clusters with a common center.

shown in Fig. 6.15; and the main part of the cluster is presented in Fig. 6.16). A shadow mapping of the grid with meshes of $250 \times 250 \mu\text{m}$ is noticeable on the detector. The edge of a magnetic pole is also well-projected (it is not shown on the figure). We note that the exposure of a detector at the same position at once after the shot in the magnetic analyzer taken from the discharge chamber did not lead to any increase in the background of nuclear tracks. We should pay attention to the fact that the slanting tracks are mainly directed downwards (for the given orientation of the figure) together with the well-pronounced dispersion centrality which is clearly seen

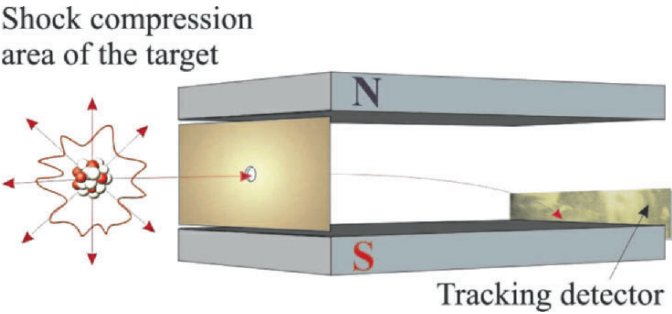


Fig. 6.15. Scheme of the experiment, in which a “giant” cluster was observed.

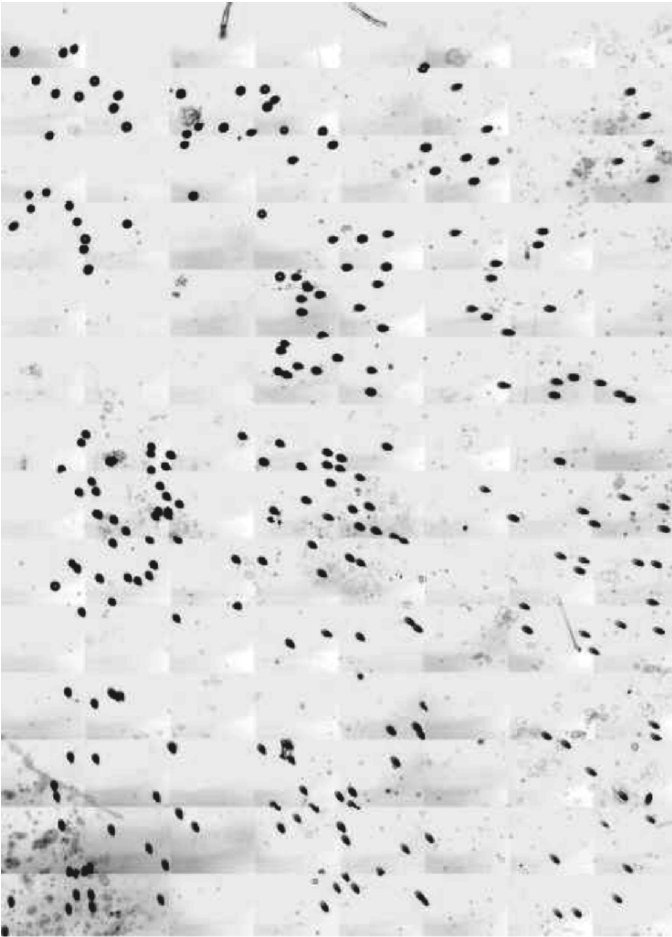


Fig. 6.16. “Giant” cluster of 276 directed tracks.

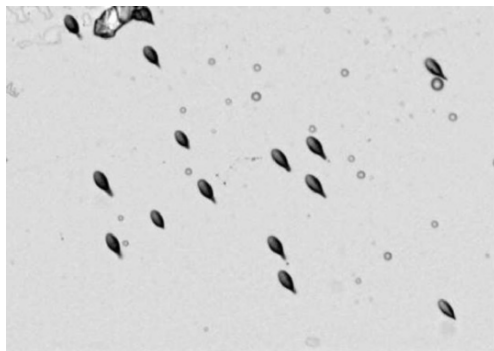


Fig. 6.17. Separate frame of the $240 \times 300 \mu\text{m}$ microscopic image shown in Fig. 6.16.

in Fig. 6.18. The character of tracks and their ideal directivity is shown in Fig. 6.17 presenting a magnified fragment of the cluster. The attempts to analyze the track cluster shown in Fig. 6.16 allow us to conclude that, like the above-discussed case, it is formed by α -particles with energies up to 8 MeV and possibly by other light nuclei. (In the mentioned cases of the observation of unique clusters, the construction of loci is hampered, because we did not realize a strict control over the etching rate of detectors in these experiments).

To overcome the doubts as for the possibility of the radioactive origin of the giant cluster (as known, the “stars” of tracks can be generated by “hot” radioactive microparticles of the type of a fixed plutonium dust particle upon a suitable exposure) and to clarify the role of the conditions of its registration, we carried out a model experiment. We put a small piece of the radioactive foil of 0.1 mm in thickness which contained Np^{237} in the magnet gap of the analyzer at the same point near the edge of the south pole of a planar magnet, where the emission center was located in the real experiment. Similarly to the real experiment, we mounted a fine-mesh metallic grid at the output from the magnet gap in front of a track detector. The registration of a radioactive cluster was performed in air at the atmospheric pressure without evacuation, which allowed us to judge on the influence of the scattering of α -particles in air on the number of arbitrary oriented tracks of the registered radioactive “star”. We took the exposure to be such that the cluster includes several hundreds of tracks.

As was expected, the simulation showed the absence of any influence of a magnetic field on the directivity of the tracks of nonscattered α -particles. The maximum energy of the α -particles from the neptunium source is 4.8 MeV, so that their cyclotron radii under conditions of the simulation reach 1 m. Therefore, even the peripheral tracks in the scope of the

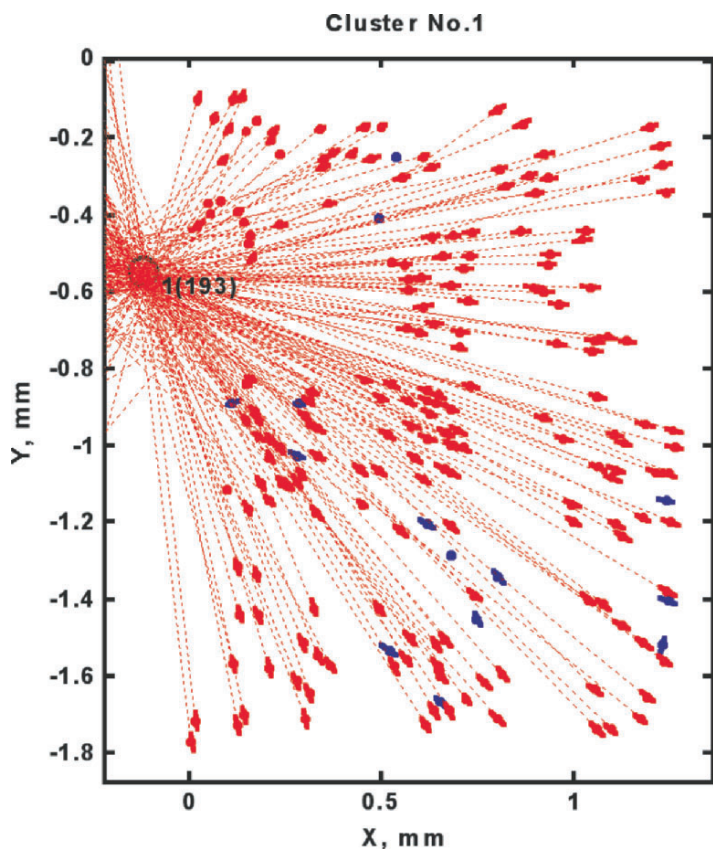


Fig. 6.18. Directivity diagram of the tracks of a cluster.

cluster are oriented to the center. Contrary to the cluster shown in Figs. 6.16 and 6.18, the model cluster turned out to have a completely isotropic angular distribution of tracks. The counting of tracks on both sides of the conditional line passing through the emission center normally to the edge of a magnet and a detector (in Figs. 6.16 and 6.18, it is a horizontal line passing through the cluster center) gave the ratio 1 : 1 for the numbers of tracks for the model “star”, whereas the cluster in Fig. 6.16 is characterized by the ratio 64 : 212. The model experiment allowed us to answer the question: To which extent does the number of scattered tracks depend on such factors as a grid on the way of particles or the presence of a gas at a certain pressure? The number of scattered tracks in the model cluster was 63% of the total number of tracks. But for the giant cluster in Fig. 6.16, this index was 11% (in Fig. 6.18, we separate 31 tracks by blue color, whose direction did not allow us to include them in the cluster of 276 tracks oriented to the center). For this reason, the model experiment gave a pattern completely different from

the “giant cluster”. Indeed, the derivation of the identically directed groups of tracks similar to those in Fig. 6.16 is impossible under conditions of the simulation. It is obvious that the elastic scattering of α -particles by atoms of air deviates considerably α -particles from the initial direction of motion already on the paths of the order of several millimeters, though the energy loss of α -particles due to the inelastic collisions occurs at a distance of at least 1 cm.

Thus, the cluster of tracks shown in Fig. 6.16 can be produced only in the scope of the vacuum cycle of one shot (about a half-hour) and cannot be referred to background manifestations of the type of a contamination of the cavity of the magnetic analyzer with a radioactive dust during the preparation of a shot. The anisotropy of the “giant cluster” can be related to the kinematics of the emitting center like the case of the cluster of 20 tracks shown in Fig. 6.13. In this case, we have to exclude the possibility of a connection of the anisotropy of the cluster directivity with the direction of the falling of a sufficiently heavy “hot” microparticle enriched by short-lived α -active nuclei, which is produced by the explosion of a target or introduced upon the preparation of the target, on a magnet. To this end, it is necessary to assume that all α -decays were realized during the time interval when the particle contacts with the magnet surface.

The pattern of a simultaneous emission of light nuclear fragments as a result of the explosion of a moving superheavy particle (of a hypernucleus?) seems to be most probable. But such an event is impossible in the framework of ideas of the orthodox nuclear physics. The second possibility consists in the attraction of a hypothetical factor of increase in the pionic background related to the collapse of a target. It is known from the physics of π -mesons that they are strongly absorbed by nuclei and remain the explosion-induced stars on track detectors which are similar to those shown in Figs. 6.12 and 6.13. In this case, the main products of the fragmentation of nuclei are α -particles. We may assume that charged π -eo are created under certain conditions in targets, where nuclear transformations occur, and can be registered at large distances from a target in the scope of their 26-ns lifetime. Small clusters of the type of those presented in Figs. 6.12 and 6.13 can be interpreted as the elementary events of pairwise pion-nucleus interactions.

However, the “giant cluster” cannot be a product of the pairwise pion-nucleus interaction, and we have to assume the presence of a pion jet or a local process of multiple generation of pions inducing the nuclear regeneration of a substance (in the given case, in the surface layer of a ferromagnet), which is equivalent, in essence, to a nuclear microexplosion. Such a hypothesis can be justified in the framework of “phantasmagoric” ideas, for example, of the possibility of the annihilation of superheavy magnetic

“grand-monopoles”, whose existence is discussed in the modern field “theory of grand unification”. The mass of heavy monopoles is sufficiently large ($>10^{13}$ GeV) in order that the S-N-annihilation generate a huge number of $\pi^\pm$ -pairs, whose paths in solids equal tenths of one millimeter. We will briefly discuss a possible scenario of the processes, which can occur in a collapsing target and can generate fragments of the nuclear matter absent in the free state primordially, in the next item of this section in terms of an alternative scheme, namely the “multiple magneto-nuclear *catalysis*”.

6.7. Discussion of the Results and Conclusions

1. The unique feature of the experimental setup IVR-1 at the Electrodynamics Laboratory “Proton-21” is its ability to create the flows of energy and plasma in the gap of a vacuum diode in order to provide the efficient destruction of a target being an element of the anode³. The explosion of the target is accompanied by the intense emission of fast protons, deuterons, and X-rays and is completed by the formation of a cumulative channel seen in the remnants of the target and a partial evaporation of the target holder and the deposition of its vapors with the changed element composition on adjacent accumulating screens.
2. The energy spectrum of fast ions measured integrally by the sorts of ions decreases monotonically as a function of the specific energy $E^* = (M/Z^2)E$ (almost linearly on the logarithmic scale [see Fig. 6.3]) and is characterized by a sharp edge at energies of the order 1 MeV. Though the “integral” spectrum of fast ions does not reveal any resonances in the region of hundreds of keV, the high-energy edge of the spectrum is characterized by some rise, which allows us to assume the presence of monoenergetic components of the beam type at the leading edge of the plasma expanding from the exploded target. Moreover, the track detectors in the magnetic analyzer and Thomson mass-spectrometer register comparatively scanty high-energy groups of protons which do not reveal themselves in Fig. 6.3 are not seen in the photo of Thomson parabolas in Fig. 6.11a. These groups are separated by some energy gap from the basic spectrum. At present, the maximum energy of the registered protons is equal to 2.7 MeV. The corresponding nonmonotonic features in the spectrum of the fastest protons registered by detectors can be associated with nuclear reactions running in a target like, for example, to the peak of 3-MeV protons in the “thermonuclear ashes” upon D-D reactions.

³This scheme is efficiently realized in both cases of the use of metallic and dielectric targets.

A similar anomaly is especially pronounced for deuterons, whose contribution is clearly seen on the $M/Z = 2$ parabola in Fig. 6.11a. As distinct from the proton-related $M/Z = 1$ parabola, the distribution of tracks over the parabola length (i.e., over the energy $E^* = [M/Z^2]E$) is nonmonotonic: the distribution of deuterons is characterized by a maximum in the region with $E^* = 0.7 \text{ MeV}$ (i.e., $E = 350 \text{ keV}$ for deuterons). In the lowest-energy region with $E^* = 300 \text{ keV}$, deuterons are absent at all, though protons on this edge of the spectrum are characterized by a monotonic increase in their number. Analogously, the tracks in the scope of the $M/Z=3$ parabola in the region with $E^* \approx 1 \text{ MeV}$, which are interpreted above as those associated with tritons and nuclei ${}^6\text{He}$, are also characterized by some growth in their density in the “head” of the parabola.

Thus, the monotonically decreasing spectrum of fast protons with energy up to 1 MeV , which are freely leaving HD, contains no information about which nuclear processes run in a solid target in the case of its successful damage. The registration of mainly the hydrogen component can be explained by that the fast ions escape from the target center at the time moment, when its shell has not been destructed. It is natural that, in this case, hydrogen nuclei pass through the target shell with the least lose of energy. In addition, heavy nuclei possessing such an energy of ions, at which the nuclei of hydrogen make deep tracks registered by us, will make no remarkable traces. However, the presence of the groups of protons with energies $> 1 \text{ MeV}$, as well as the anomalous additional beams of heavy hydrogen and helium, testifies to the running of the processes of nuclear transformations at HD.

3. The total number of fast protons and deuterons estimated by the integration of this spectrum is of the order of 10^{14} particles. It is a very high value if we take into account the comparatively low energy content of our setup (about 3 kJ) and especially of the driver proper (up to 300 J) as compared to the megajoule scale of experiments with the plasma focusing which yield about 10^{15} “fast” ions (see Ref. 202). The great number of particles derived in our experiments cannot be caused by the hydrogen contamination of a target (the surface desorption by hydrocarbons and water, which gave the yield of 10^{10} protons in Ref. 203, and the hydrogen dissolved in the target can ensure at most 10^{12} protons in our experiments). The excess level of emission of protons over the initial contamination by the order of two to three and the still more essential excess of deuterons (we consider only track-forming particles) also testify to the presence of the processes of nuclear transformations.

4. The track registration of deuterons, being the rightful (along with protons) component of the emission of HD, is one more proof of the presence of anomalous nuclear processes at HD in our experiments. Strictly saying, this conclusion is valid only for the region of paths, in which the track loci reveal the splitting and demonstrate the presence of deuterons. The identification of particles is realized in the region of paths determined by the TDS method, $R = 4$ to $12\ \mu\text{m}$, on detectors CR-39 exposed in the obscure-chamber and magnetic analyzer (Figs. 6.9 and 6.10B, respectively), as well as in a Thomson mass-spectrometer (Fig. 6.11b). In this case, we may consider that the loci, e.g., in Fig. 6.10B, represent proportionally protons and deuterons from the stripes equivalent by the density of tracks and the width on a detector in the magnetic analyzer. By restricting ourselves by the corresponding intervals of real paths and energies and basing on the energy spectrum shown in Fig. 6.3, where the identified deuterons demonstrate large values of the cyclotron radius satisfying the condition $E^* > 0.6\ \text{MeV}$, we get a rough lower estimate for the number of deuterons to be of the order of 10^4 , whereas the integral number of tracks associated with this spectrum is 10^6 . This underestimated evaluation corresponds to the content of at least 1% of deuterons in the total yield of fast ions under conditions of the absence of a preliminary deuteration of the target material. This estimate is confirmed upon the quantitative analysis of the track signal of deuterons which is clearly pronounced on the corresponding parabola in Fig. 6.11a.
5. The registration of the products of nuclear reactions as the tracks of fast deuterons or the clusters of tracks in the corpuscular emission of HD requires to consider them in one context with other data on the registration of heavy nuclei in the decay products of a damaged target on the setups of the Electrodynamics Laboratory "Proton-21". Possible conceptual alternatives are analyzed in the separate chapter. Here, we briefly consider the concepts of "multiple magneto-nuclear catalysis", whose ideas allow one to understand, in particular, why the target explosion products contain a lot of protons and deuterons, rather than, for example, α -particles.

The idea consists in that the action of a very strong magnetic field (whose intensity is unknown, but is presumably attainable under laboratory conditions) induces the dissociation of "colorless" aggregates of quarks, hadrons, at which the superstrong (not decreasing with increase in the distance, according to chromodynamics) interaction is realized as the repulsion of quarks, whose spins and magnetic moments are oriented with the violation of the ordinary rules valid

for strongly bound hadronic states. This effect can be named as the “color disconfinement in a strong magnetic field”. The process of scattering is running through a bound, *virtual* compound nucleus in the form of a cloud of the nuclear matter (a quark-gluon-electron plasmoid) which has a density much less than the nuclear one and is spatially expanded over the interatomic distance. After the switching-off of the magnetic field, this cloud returns to the state of ordinary strongly bound hadrons, mesons, and baryons, with steadily ordered spins and magnetic moments. The composition of nuclei in the final state can differ from the initial one. In this case, the creation of protons and deuterons is more preferable as compared to α -particles, since the appearance of even-even nuclei with zero magnetic moment is obviously improbable in a superstrong magnetic field.

The above-presented assumption about the possibility for a superstrong magnetic field to influence the intranuclear processes can be valid for weak nuclear interactions defining the stability of nuclei relative to β -decays. At present, we can consider the defining influence of the electron system of atoms on the natural β -stability of nuclei as the established fact (see Ref. 206), which is equivalent to the account of effects of the collective electromagnetic field of electron shells on the nucleus at the center of an atom. Because the β -activity redistributes nuclei in the limits of isobaric families, we cannot explain the appearance of new light or heavy nuclei by the action of a superstrong magnetic field on the weak interaction of nucleons, but can understand why the collective magneto-nuclear catalysis leads to the observed effect of the β -stable composition of synthesized nuclei. The absence of medium and heavy β -active nuclei together with stable representatives of the isobaric families in products of the explosion-damaged target can be explained by the strengthening of the weak interaction and the increase in the rate of β -decays else in the limits of the action of a superstrong magnetic field (on the scale of forces of the weak interaction), when the insulating nuclei realize their β -activity more rapidly, so that the releasing nuclei turn out to be absolutely stable. Nuclei of tritium and ${}^6\text{He}$ observed in some experiments with the help of a Thomson analyzer are apparently referred to manifestations of the activating action of fast protons and deuterons on stable nuclei of the slower component of dispersing products of the exploded target. On the last stage of the explosion, acts of the binary nuclear interaction described in the framework of orthodox nuclear physics can be realized.

A development of the model of generation of a superstrong electromagnetic field by HD is possible in the framework of ideas of the “Hall” or “electron” MHD theory (see Ref. 204), but is not the aim

of this work. We can perform the relevant estimates in order to show the possibility of a modification of the composition of a quite large amount of the target material and the anode plasma with regard to the energy content of discharges on the setup at the Electrodynamics Laboratory "Proton-21". But such estimates would contain a significant degree of arbitrariness in our considerations. As for the idea of "multiple magneto-nuclear catalysis" possibly presented above for the first time, we can indicate no references in the literature devoted to the possibility of a realization of this mechanism in really attainable electromagnetic fields (except for the last reference to Ref. 206 concerning the dependence of the weak nuclear interaction on a state of the electron atomic subsystem). Only further studies, the development of computational models, and the definitive formation of chromodynamics can clarify the problem posed above.

EXPERIMENTS ON THE NEUTRALIZATION OF RADIOACTIVITY

S. V. Adamenko and A. S. Adamenko

Since the year 2000, researchers at the Electrodynamics Laboratory “Proton-21” have pursued experiments on the artificial initiation of the collapse of a substance in solid targets. In this case, a target is destroyed by the explosion from the inside. A significant part of target explosion products is deposited on an accumulating screen located in the close vicinity of the target during the experiment.

The studies of the element and isotope compositions of products of the explosion-induced destruction of targets and the analysis of possible reasons and mechanisms of their appearance have gradually led us to the conclusion that, in the process of collapse, there occurs a nuclear transformation of the various chemically pure materials of targets with the formation of chemical elements in a wide range. In this case, the formed elements and their isotopes contain no radioactive ones.

The absence of radioactive isotopes in products of the nuclear transformation of targets gives evidence for that the process of collapse is associated with the running of collective reactions which result in the formation of isotopes most favorable by energy (by the specific binding energy per nucleon). In this case, if we use the targets made of radioactive materials, we may expect the origination of stable isotopes of various chemical elements in the collapsing region.

This conclusion is based on a number of the following logically connected assumptions. The process realized in the Electrodynamics Laboratory “Proton-21” is not reduced to a simple transformation of some nuclei into the others by means of the direct fusion or fission, i.e., it does not bear an individual character. It is a basically collective process passing through the intermediate stage of the formation of a very large electron-nucleus cluster. Such a cluster can emit other nuclei allowed by the relevant conservation laws during its growth or after the termination of this process. Nuclear transformations of isotopes occur in our experiments, most likely, by passing through the stage of collective and strongly correlated states leading to the complete collapse.

Due to a strong compression of the active region, nuclei lose their individuality in a volume being collapsed. We expect the formation of a giant macronucleus or, to be more exact, an electron-nucleus macrocluster which undergoes a number of global transformations and begins to decay with the emission of new nuclei. Without complications, we may imagine that, on the first stage of transformations, all interacting nuclei are “decomposed” into separate nucleons with the simultaneous formation of a giant electron-nucleus macrocluster. Then, other stable daughter nuclei are formed from these nucleons and leave the volume of the parent macrocluster. Apparently, the specific mechanism of formation of nuclei in the macrocluster volume is similar to the process of α -decay of ordinary heavy nuclei. First, a new nucleus is accidentally formed in the cluster volume. If this nucleus is sufficiently stable, it can survive, not decay in the cluster volume, and escape outwards (of course, if such an escape is energy-gained). It is obvious that such nuclei are nonradioactive. Indeed, radioactive nuclei are those, in which the proton-neutron composition is nonoptimum or protons or neutrons are in a nonoptimum state. But any accidentally formed nucleus with nonoptimum composition is unstable and is easily destroyed in the cluster volume. Only stable nonradioactive nuclei have the highest probability to survive in the cluster volume and to leave its volume. In this case, it is not significant whether the initial nuclei are radioactive or stable.

It is well known that the problem of the utilization of radioactive wastes, being quite urgent for many developed countries, remains unsolved. On the other hand, the transformations of the materials of targets, being a result of the experiments performed at the Electrodynamics Laboratory “Proton-21”, indicate the potential possibility for the neutralization of radioactive wastes.

With the purpose to practically verify such a possibility, about 100 experiments on the neutralization of a radioactive material artificially introduced into a target were carried out at the Electrodynamics Laboratory “Proton-21”. To conduct these experiments we had to solve a technological task related to the fact that only the central part of the target volume is regenerated during the explosion. Thus, it was necessary, prior to experiments, to set a microscopic core made of a radioactive material at the center of a nonradioactive shell. In this case, it was important to provide the ideal contact between the shell and the core. In our experiments, the requirements to a core material were limited by several points: it should be activated by thermal neutrons, have a great half-life period, have high-energy lines in the spectrum of γ -emission, and admit an easy mechanical processing. Under the same conditions, the shell material must not be activated or should reveal a rapidly decaying induced radioactivity. The longitudinal section of a target

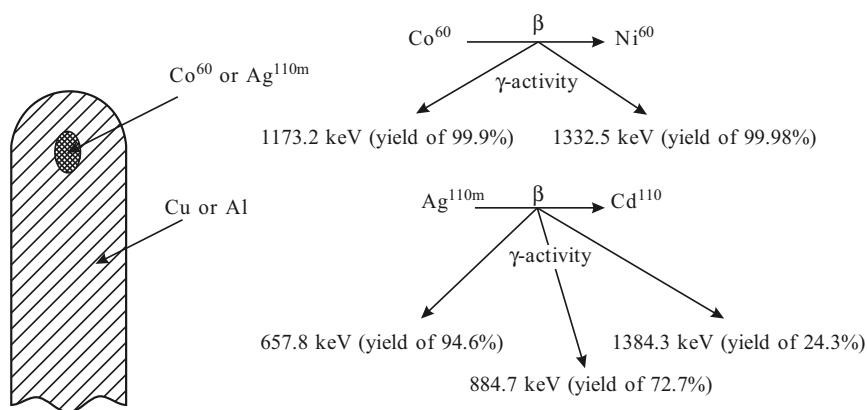


Fig. 7.1. Scheme of a target and the main lines of γ -emission of the radioactive isotopes used in targets.

used in the experiments on the neutralization of radioactivity is schematically presented in Fig. 7.1. There, we also give the schemes of natural decay of the used radioactive isotopes and the most probable energies of γ -quanta emitted in these decays. As the core material, we took cobalt or silver activated with neutrons. The energy-concentrating shell of a target was made of copper or aluminum.

The preliminary irradiation of targets with neutrons (activation) resulted in the formation of isotopes Co^{60} in a core made of Co by the scheme $\text{Co}^{59} + n^1 = \text{Co}^{60}$ and Ag^{110m} in a silver core by the scheme $\text{Ag}^{109} + n^1 = \text{Ag}^{110m}$. The concentrations of radioactive isotopes in activated cores were about 0.001% of the total amount of nuclei in the capsule. Such a degree of activation was chosen in view of the requirements of radiation safety. In the first case, the variations in radioactivity were detected by two lines of the γ -emission of Co^{60} : 1173.2 and 1332.5 keV. In the second case, we used three most intense lines of the γ -emission of Ag^{110m} : 657.8 keV (yield of 94.6%), 884.7 keV (yield of 72.7%), and 1384.3 keV (yield of 24.3%).

Experiments were performed using a hermetic cylindrical caprolan chamber. The shell with a core was positioned at the center of this chamber. The measurement of radioactivity prior to the execution of an experiment was realized for one of the two different distances from the center chamber to the surface of the detector (75 mm or 755 mm depending on the radioactivity of a target) or for both distances. The measurements on some targets cannot be carried out at the greater distance because of their weak activity and were realized only at a distance of 75 mm. The measurements were carried out with a spectrometric germanium-lithium γ -detector DGDК-100 V-3.

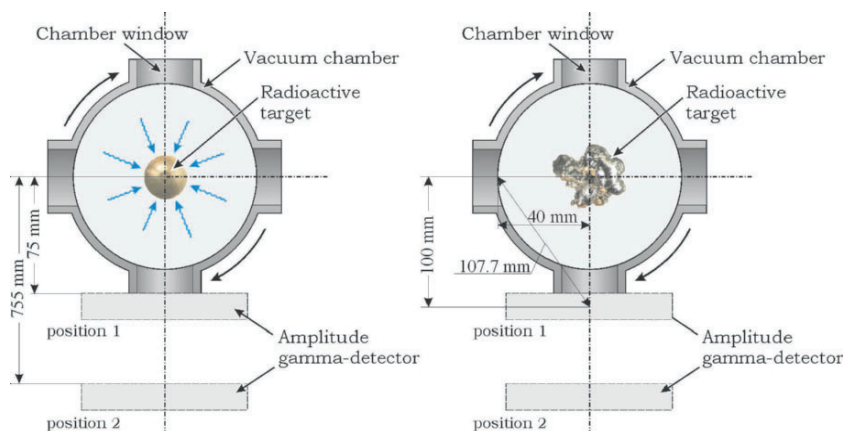


Fig. 7.2. Scheme of the measurement of the chamber radioactivity with a target before and after the execution of an experiment.

In Fig. 7.2, we schematically display the chamber with a target placed on the detector. At each distance, the radioactivity was measured in turn at four fixed positions of the chamber. A change of the position was realized by the rotation of the chamber around its symmetry axis by 90° . As the initial radioactivity for a given distance, we took the arithmetic mean value by four measurements. The chamber with the open valve was mounted in the vacuum chamber of the setup intended for the high-energy action and was vacuumized. The valve was closed several seconds prior to the beginning of the experiment, and the chamber became hermetic. After the high-energy action on a target by the electron beam, the chamber remained hermetic up to the completion of all measurements. The measurements of radioactivity inside the chamber after the experiment were performed at the same distances and in the same positions as prior to the experiment.

Due to the high-energy action, the target explodes, and its remainders are sprayed over the inner surface of the chamber. We indicate the following two factors as the sole sources of errors upon the measurement of radioactivity after the completion of the experiment: the change of the effective distance from the radioactive material (on the average by each separate fragment of the target) to the detector surface and the change of the absorption of γ -quanta by the chamber walls due to the redistribution of the radioactive material in the chamber.

Consider the first factor. The movement of the radioactive material in parallel to the detector surface within the inner space of the measured chamber (at most 40 mm from the center) does not practically affect the counting rate (even in the case of measurements at the smaller distance),

because the detector crystal has a cylindrical form (53 mm in diameter and height) and absorbs γ -quanta by the lateral surface as well. In this case, the distance from the radioactive material to the center of the detector crystal at the maximum distance from the chamber center can vary by at most 7.7%. Now consider the change in the counting rate upon the movement of the radioactive material in the plane perpendicular to the detector surface by using a simple example. We recall that, at the beginning, the radioactive material was at a distance of 75 mm from the detector surface (about 100 mm to the center of the detector crystal) (see Fig. 7.2). Let its radioactivity be equal to X . After the experiment, the radioactive material can move at most at the distance equal to the inner radius of the chamber (40 mm). By the subsequent washing away of the sprayed radioactive material from the various parts of the inner surface of the reaction chamber, we established that the majority of the radioactive material was distributed near the plane that passes through the target center and is perpendicular to the chamber axis. In Fig. 7.2, we display the cross-section of the chamber just in this plane. We now estimate the effect of the process of redistribution of a radioactive material on the registration of the flux of γ -quanta. Let us divide the radioactive material into four equal parts. We remove one part from the detector by 40 mm, bring nearer to the detector the second one by the same distance, and move the third and fourth parts along the detector surface in opposite directions by 40 mm from the center (see Fig. 7.2). The distances from the third and fourth parts to the detector center become equal to 107.7 mm. The counting rate of the detector (K) is inversely proportional to the square of the distance from the radioactive material to the detector. Hence, the counting rate $K_1 = 10^{-4}X$ before the division of the radioactive material and, after it, $K_2 = 0.25 \cdot X/60^2 + 0.25 \cdot X/140^2 + 0.5 \cdot X/107.7^2 = 1.25 \cdot 10^{-4}X$.

Thus, after the redistribution of the radioactive material within the chamber without any change in its amount, the mean counting rate of the detector can be increased by 25%. In view of the fact that some part of the radioactive material can move along the axis of the cylindrical chamber (in the limits of 30 mm from the place of its initial position), the derived value will be somewhat less (approximately 10% to 15%). The same numbers were derived experimentally upon the intentional arrangement of a material with known radioactivity at different points of the chamber and the comparison of the derived results of measurements with those in the case of the radioactive material located at the chamber center. While proving this fact, it is worth noting the results of measurements of the chamber upon the use of target No. 5563/42 (experiment No. 4372) (see Table 7.1). According to the results of measurements at a distance of 75 mm, we get the increase in

radioactivity (by 16.5% on the average), which is related to the geometry of the dispersion of a radioactive material. But by the results of measurements on the same chamber at a distance of 755 mm, no change in radioactivity was practically observed.

Hence, we can conclude that, upon the measurement of the chamber radioactivity at a distance of 75 mm, the percentage of neutralization will be, as a rule, underestimated. Upon the departure from the detector, the chamber becomes similar to a point source, and the geometry of a redistribution of the radioactive material inside the chamber will not affect the counting rate of the detector.

This fact is confirmed by the results given in Table 7.1. The data show that, upon the measurement of the radioactivity of bodies at a distance of 755 mm, the percentage of a decrease in radioactivity is higher practically

Table 7.1. Changes of the chamber radioactivity in experiments with silver targets.

Expe- riment	Radioactivity decrease			
	657-keV line	884-keV line	1384-keV line	Mean
3902	7.9	7.1	1.6	5.5
3902*	10.8	11.3	13.3	11.8
3903	5.5	3.9	2.7	4.0
3903*	11.0	10.6	10.9	10.8
3919	8.5	5.2	3.8	5.9
3919*	9.8	11.5	9	10.1
3925	9.2	8.1	8.6	8.6
3925*	13.7	13.8	15.1	14.2
3936	8.0	7.2	10.1	8.5
3936*	10.6	10.2	10.1	10.3
4362	6.1	3.9	-2.4	2.5
4362*	9.8	10.6	10.0	10.1
4372	-15.9	-16.0	-17.7	-16.5
4372*	-2.3	-3.9	1.2	-1.7
4836	6.7	5.9	2.5	5.0
4836*	14.1	16.8	14.2	15.0

* experiments, in which the measurements of radioactivity were performed at a distance of 755 mm.

in all the experiments than that upon the measurement of the same bodies at a distance of 75 mm.

As for the second factor, it would be controlled upon the execution of experiments with silver targets. The results of experiments on the neutralization of the radioactivity of silver targets are given in Table 7.1. The used isotope of Ag has a wide spectrum of γ -emission from 657 to 1384 keV. Thus, in the case of an increase in the absorption, γ -quanta with lesser energy will be absorbed stronger, and the decrease in radioactivity by the 657-keV line will be considerably greater, than that by the 1384-keV line. Such phenomena were sometimes observed upon the measurements of bodies at a distance of 75 mm. However, upon the measurements of bodies at a distance of 755 mm, the decrease in radioactivity was practically the same by all the lines. Hence, upon the measurement of bodies at a distance of 755 mm, no change in the absorption of γ -quanta occurs because of the redistribution of a radioactive material over the inner surface of the chamber. Thus, upon the measurements at the greater distance, both above-described factors, as a rule, do not influence the results of measurements, and the values of radioactivity of the chamber before and after the experiment that were calculated by the results of four measurements are quite exact.

Upon the measurements of variations in the radioactivity of bodies at a distance of 75 mm, two factors contradict each other: on the one hand, the percent of neutralization is underestimated at the expense of the change in the distance from microfragments of the activated capsule to the detector surface. On the other hand, it can be overestimated in separate experiments at the expense of the increase in the absorption of γ -quanta by the chamber walls. The experiments, in which the measurements were performed at both distances, show that the percent of neutralized radioactivity derived with the use of the results of measurements at the smaller distance is usually underestimated. The experiments, in which the correction on a change in the absorption of γ -quanta should be made, can be always distinguished by different percents of the reduction of radioactivity for different lines of the γ -spectrum.

The results of measurements of the radioactivity of the chamber before and after experiments showed that, upon the successful execution of an experiment (under the correct operation of the setup electronics), radioactivity is reduced in the limits from 10% to 38%. The results of these experiments with the use of different targets are given in Table 7.2.

At present, the derivation of values higher than those given in Table 7.2 is limited, in the first turn, by the circumstance that only a part of the core volume with a radioactive material is involved in the process of

Table 7.2. Results of the experiments on the neutralization of radioactivity.

Experiment	%	Experiment	%	Experiment	%
2397	37.8	2676	11.0	3277*	16.1
2481	10.0	2728	18.5	3279*	17.7
2534	22.7	2743	14.7	3902*	11.8
2560	11.9	2770	25.0	3925*	14.2
2582	9.1	2778	22.2	3936*	10.3
2588	36.6	2781	12.3	4097*	6.7
2600	20.5	2782	13.0	4362*	10.1
2673	37.4	2785	33.7	4836*	15.0

* experiments, in which the measurements of radioactivity were performed at a distance of 755 mm.

collapse and regeneration. This depends on the following temporary technological factors:

- Microscopic sizes of the core did not allow us to manufacture it with ideal form, namely, in the form of a cylinder with rounded corners at its bases.
- The presence of the boundary between the core and the shell reduces considerably the volume of a core domain involved in the process of collapse.
- The uncertainty in the location of the point of collapse on the target axis, which cannot be predicted at present and depends on the initial parameters of the discharge that is only partially controlled, does not allow us to place the geometric center of a capsule at the needed point of the inner space of the shell prior to the experiment.

ISOTOPE AND ELEMENT COMPOSITIONS OF TARGET EXPLOSION PRODUCTS

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Shock compression of solid targets in a vacuum chamber experimental setup (residual pressure $\sim 1 \times 10^{-3}$ Pa) was achieved with the help of an electron beam, used as the primary carrier of the concentrated energy (up to 1 kJ of energy for the impulse time of the order of 10 ns). This beam caused the creation of a superdense state of matter with $\rho \gg 10^{26}$ nuclei/cm⁻³ in a microvolume of the targets in the vacuum chamber. In all the cases studied, the power density in the domain of compression, by various estimations, exceeds 10^{22} W · cm⁻³. As a result of this shock compression, targets were destroyed by the explosions which evidence shows came from the inside of the specimens. The indicated compression process is usually accompanied by the radial dispersion of the target substance with evidence of its deposition on the body walls of a vacuum chamber of the experimental setup and on a special accumulating screen upon which the targets rest. The accumulating screens have the form of a disk about 15 mm in diameter and 0.5 mm in thickness. A part of explosion products were released in the form of a gas phase into the residual atmosphere of a vacuum chamber of the setup.

Following with the above discussion, to obtain specimens of the products of the explosions of the targets, we took gaseous samples from the reaction chamber after the high-energy action on a target. We sampled layers of the condensate deposited on the vacuum chamber walls. We also examined material scattered on the accumulating screens after the explosion, as well as, sampled the reaction products on the surfaces of the remnants of the target and on the craters that were formed in the targets in the form of an axial channel after the explosions. We studied the isotopic and element compositions. We emphasize the following circumstances. In the case of specimens obtained from fragments found on the walls of the vacuum chamber, we dealt usually with the condensate formed after several hundreds or even thousands of experiments. But in the case of the specimens taken from the targets and accumulating screens, we were able to study the products of individual explosions. For these reasons, as specimens for

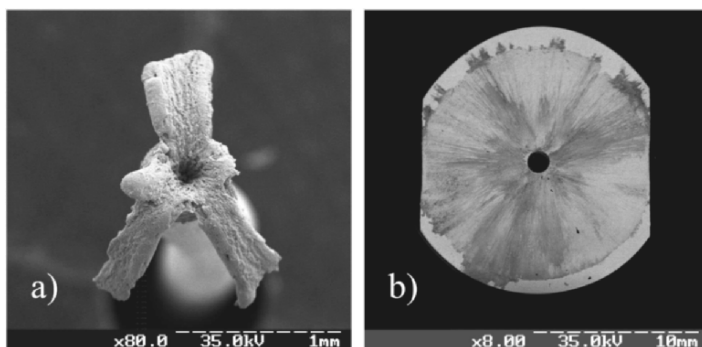


Fig. 8.1. Scanning electron micrographs (the secondary electron image mode) of a typical target after the explosion (a) and a typical accumulating screen with deposited products of the explosion (b).

the study of explosion products, we mainly used the exploded targets and accumulating screens after each explosion. In Fig. 8.1, we present the micrographs of a typical target after the explosion (a) and a typical accumulating screen with explosion products deposited on it (b) derived on a scanning electron microscope in the secondary electron image mode.

The analytic studies of the products found after the explosive destruction of targets were preceded by the study of both dynamic parameters of the process of expansion of a plasma bunch arising in the zone of shock compression and the spectral parameters of its emission by detector methods (see Chapters 4–6). The emission spectrum of a plasma flash initiated by shock compression of the target substance was registered in the optical, X-ray, and γ -ranges. We also made observations of tracks of nuclear particles present in the plasma of the dispersed substance of destroyed targets. These observations give evidence of the existence of intense artificially initiated processes of nucleosynthesis and transmutation in a microvolume of the target substance which has undergone compression up to superhigh densities. The last circumstance has caused us to engage in a course of extensive analytic studies of the products that result from these newly discovered processes of deep transformation of target substances that occur during their compression to superdense states. By virtue of the fact that new chemical elements originate in these nuclear reactions and their isotope composition can differ from the normal natural proportions, we found the need to study the element and isotope compositions of explosion products by using known analytic methods. It is obvious that the establishment of the appearance of chemical elements in the explosion products that were earlier absent in the composition of the initial materials of targets and accumulating screens (structural details of the experimental chamber that participate in the process of

explosion), as well as the registration of anomalies of their isotope composition, would uniquely testify that the target explosion products contain nucleosynthesis products. However, we note that, up to now, science has no answer to the following questions: What is the ratio of isotopes in the products of a natural explosive nucleosynthesis?; What factors does that ratio depend on?; In what bounds will it vary in dependence on the parameters of the corresponding explosion?; What basic schemes of formation of the ratios of the concentrations of isotopes exist in the Nature (Refs. 70–73)? It could be that some of the processes we observe are mimicking the processes which formed the elements found in the Nature. Therefore, the very fact of the preservation of natural proportions in the ratio of the concentrations of the stable isotopes of chemical elements observed in the specimens can serve neither the proof of their belonging to natural nuclides, nor the disproof of their artificial origin.

8.1. Isotope Composition of Explosion Products

Following our hypothesis that we observe massive nucleosynthesis in our experiments, a first step would be to compare the isotope distributions of the experimental results with those which have been established for natural occurring substances. In this case, of importance is the established abundance of the elemental isotopes. In contemporary science, the mean abundance of nuclides is determined by the totality of the data of geochemistry, space-chemistry, and astrophysics by studying the composition of specimens containing the earth's elements, meteoritic and lunar substances. Also the electromagnetic emission spectra of the sun, stars, and the interstellar medium, as well as by determining the content of nuclides in the solar corpuscular emission and galactic cosmic rays (see Ref. 71) serve as useful data to establish natural isotopic abundances. At present, the isotope composition of all known substances have been studied succinctly and well only for the sun's system, and the most reliable data are obtained from meteorites (see Ref. 74). Therefore, scientists consider the content of nuclides in meteorites as a standard upon which they base the systematization of their abundance in the Nature. The systematization of these data was carried out by A. Cameron in 1982. The data on the abundance of nuclides are usually normalized on the content of Si and are presented in the form of a curve on the logarithmic scale, the so-called standard curve of the abundance of nuclides. Now the mentioned data are the main experimental basis of the theory of nucleosynthesis in the Nature (see Ref. 75).

It is surprising that, despite the diversity of the processes forming the substances of the Universe, the compositions of the majority of stars, galaxies, and the interstellar medium fall mainly on the standard curve of

the abundance of nuclides. Some well-known deviations from this curve are caused by various physical reasons understandable in many respects. However, the stars with anomalous chemical composition constitute only about 10% of the stars positioned near the main sequence (the Hertzsprung–Russell diagram) (see Ref. 71). The above-written yields that if the processes of nuclear regeneration of a substance running closely to one of the natural scenarios are realized under conditions being satisfied upon the explosions of targets in the experiments performed at the Electrodynamics Laboratory “Proton-21”, it would be rash to assume that the abundance of nuclides in the products of the laboratory nucleosynthesis should be basically different from the standard one. We also note that the creation of a substance in stars from the primary gas containing 75 mass. % of hydrogen and 25 mass. % of helium gives practically the results identical to those observed after the regeneration of a substance in stars of older generations formed from the interstellar gas being several times involved in the star-forming cycles and containing the whole collection of heavy chemical elements. That is, as if the result does not depend on the starting “fuel” (Refs. 71, 76). For example, the sun is considered to be a star of the third generation formed about 5 billion years ago (see Ref. 76), and the curve of the standard abundance of nuclides, which corresponds to the overwhelming majority of objects in the Universe, is constructed on the basis of the data on the composition of the meteoritic substance belonging to the solar system and, as a minimum, already twice being in the interior of parent stars.

However, based on the above-presented discussion, we should not conclude that the search for the anomalies of the isotope composition of the products of a laboratory nucleosynthesis is a hopeless study. First of all, this is related to the idea that our investigation should be somewhat different than what standard theories would lead us to. This is caused by our observation that the earth’s substances formerly believed to be created in stars appear to be in a state of evolution. In them, the processes of slow radioactive decay and migration of chemical elements (their concentration, scattering, isotope exchange, etc.) are continuous. For the earth existence period, these processes have led to a significant differentiation of its substances (see Ref. 77) which is observed to a high degree in the place where there is a jump of the density distribution of its substances. That is, this occurs at its boundary: in the narrow layer of the earth’s crust, including the atmosphere and the hydrosphere. Here, the distribution of nuclides has a number of specific peculiarities conditioned by the above-mentioned processes, by which it differs from the standard abundance of nuclides characteristic of the substance created in the natural nucleosynthesis. Therefore, the search for anomalies of the isotope composition of target explosion

products must be mainly focused along the directions, where the abundance of nuclides in the earth's crust is different from the standard one. In this case, it is obvious that such anomalies cannot be a consequence of any natural contaminations, because the latter could only decrease their magnitude.

8.1.1. Isotope Composition of the Residual Atmosphere of the Vacuum Chamber

We can observe many bright examples of the profound differentiation of the earth's substance in its atmosphere. Here, the differentiation is mostly a consequence of two continuous processes: a selective loss and a selective replenishment of the atmosphere (see Ref. 77). The atoms of certain gases leave the atmosphere as a result of their interaction with the sun's corpuscular emission, and the intensity of this process increases with decrease in the mass of a nuclide. This leads to the shifts of isotope ratios for atmospheric gases to the benefit of nuclides with greater masses. It is for this reason that atmospheric hydrogen is noticeably enriched by deuterium. In atmospheric hydrogen, the ratio D:H=1:4500, whereas it is equal to about 1:6000 (Refs. 78–81) in ordinary water. The mentioned anomaly is small by virtue of the fact that the losses of atmospheric hydrogen are permanently replenished by hydrogen with the normal distribution of nuclides due to the photochemical decomposition of water vapors. An analogous situation is observed for oxygen and nitrogen, whose losses are completely compensated by their influx. The supply of nitrogen occurs due to the decay of nitrous compounds (Refs. 78, 82).

A quite different situation is observed in the case of inert gases. They are not supplied to the atmosphere in the process of degassing of the earth's crust in amounts such that their losses be compensated. As a result, they are characterized by a considerably lower content in the earth's atmosphere than in space. Moreover, the shift of their isotope ratios to the side of nuclides with greater masses noticeably grew during certain geological epochs (Refs. 77, 78, 83). The above is mainly referred to neon, krypton, and xenon, whose losses are practically not compensated (Refs. 84–92). As for helium and argon, the situation under consideration looks much sharper (Refs. 78, 93–98): their influx from the earth's crust to the atmosphere bears exclusively a selective character and deviates their isotope ratios from the standard ones to a still greater extent. Both go in the atmosphere in the form of the most massive nuclides: helium appears practically completely as nuclide ^4He created in the α -decay of natural radionuclides (mainly Th and U), and argon appears only in the form of nuclide ^{40}Ar created in the radioactive decay of nuclide ^{40}K by "K-capture".

As a result, the share of nuclide ^3He in atmospheric helium is approximately 1×10^2 times less than that in cosmic helium. However, this anomaly is not so noticeable by virtue of the fact that ^3He is most rarely met in the Nature among the stable nuclides, and its share in both cases is extremely low. In atmospheric helium, it is equal to about $1 \times 10^{-4}\%$. Prior to the discussion of the anomalies of the isotope composition of atmospheric argon, we note that the basic nuclides of all light chemical elements up to calcium are stable nuclides with minimum masses. The exception is ^4He due to the extraordinary high-energetic stability of its nucleus. In agreement with this rule, cosmic argon is represented by nuclides ^{36}Ar and ^{38}Ar , whereas ^{40}Ar is practically absent in the Universe (see Ref. 73). As for atmospheric argon, the anomalies of its isotope composition are significant because of the mentioned reasons (Refs. 78, 96–98). In atmospheric argon, the basic nuclide is ^{40}Ar ; its share is about 99.6%, and it originated practically wholly from ^{40}K . This fact explains one of the anomalies of the Periodic table. Contrary to the initial principle of its construction, which involves the pure use of atomic masses, argon is positioned in the Periodic table before potassium. If light nuclides were dominant in argon such as in the neighbor elements (as it occurs in the elements in space), then the atomic mass of argon were by 2 to 3 units less, and there would be no anomaly present in the Periodic table.

Thus, the above-presented information yields a preliminary conclusion that if inert gases are created upon the explosions of targets, then the abundance of their nuclides would likely differ from those found in the atmosphere, because the latter arose due to the selective migration of the nucleosynthesis products rather than due to the immediate nucleosynthesis. The isotope composition of inert gases should be analyzed taking gas samples from the residual atmosphere of the reaction vacuum chamber of the experimental setup. Indeed, it is practically impossible to register them in solid products of the explosion of a target, which are deposited on an accumulating screen, because of their low reactive ability and solvability (Refs. 78, 83). Moreover, the analysis of gas samples as compared to that of solid specimens has a number of very important advantages. First of all, gas samples have a homogeneous composition as distinct from solid products of explosions and contain a small number of chemical elements. These circumstances are extremely important for both the registration of mass-spectra being simple in their interpretation (a small number of mass-peaks and the absence of their superpositions) and the derivation of reliable results of the analysis.

The studies of the isotope composition of inert gases were performed at the Institute of Geochemistry, Mineralogy, and Ore-Formation (IGMO)

of the National Academy of Sciences of Ukraine specializing in the solution of problems of the determination of the age of rocks by the isotope ratio $^{40}\text{Ar}/^{36}\text{Ar}$. Main problems of the investigation were to determine the isotope composition of argon and to establish the share content in gas samples of other inert gases. The analysis of gas samples was carried out with the use of a mass-spectrometer MI-120IG intended for the determination of the isotope composition of microamounts of inert gases. A source of ions in the mass-spectrometer was a plasma formed as a result of the electron ionization of the analyzed gas presented in a measuring chamber. The separation of ions in the mass-spectrometer was performed with the use of a magnetic analyzer of masses (see Ref. 99).

For the investigation, samples of a gas were taken directly from a pressurized reaction chamber (its volume $\sim 0.5\text{ dm}^3$) after the explosion of a target with the help of a special unit, using a vessel made of molybdenum glass (its volume was 0.5 dm^3) as a sampler and supplied with valves DU-6 with aluminum seals and a hollow needle with an inner diameter of about 3 mm. The reaction chamber and sampler were preliminarily evacuated through the common main of a vacuum system by a diffusion oil-vapor pump up to a residual pressure of less than 2×10^{-4} Torr. Directly prior to the experiment, their volumes were cut from the common vacuum system with the help of valves. Then the target was exploded, leading to an increase in the pressure in the reaction chamber usually above 1×10^{-3} Torr due to off-gassing from the specimens. The sampling of a gas was realized in approximately 30 s after the explosion by puncturing a rubber seal of the reaction chamber with a hollow needle of the sampler. Through the needle, the created and residual gases flow to its volume due to the difference of pressures. The duration of a sampling did not exceed 10 to 15 s.

Prior to the execution of measurements, the taken gas mixture was subjected to the purification from residual gases. The first stage of purification consisted in passing the initial gas mixture through Ti and CuO, activated and heated to 923 K. Then the gas mixture (the mixture of inert gases) was absorbed by activated charcoal and a cryogenic trap at a temperature of liquid nitrogen. In the second stage of purification, the mixture was again passed over Ti heated to 723 K. Then the purified mixture of inert gases was fed to the measuring chamber of the mass-spectrometer, the time moment of the injection being fixed. All the further measurements were performed in the static mode, i.e., the measuring chamber of the mass-spectrometer closed. Its construction was such that the preliminarily established rarefaction could be held in it for 2 to 4 hours in the isolated mode at a level not less than 1×10^{-7} Torr. The ion currents of the analyzed nuclides were derived by means of their extrapolation (linear or exponential) to the injection

Table 8.1. Ratios of the basic isotopes of Ar and Kr in the reaction chamber atmosphere after the explosion of a target.

Sample	$I(^{40}\text{Ar})$, nA	$I(^{36}\text{Ar})$, nA	$\frac{I(^{40}\text{Ar})}{I(^{36}\text{Ar})}$	$\delta^\dagger$, %	$\frac{I(^{84}\text{Kr})}{I(^{86}\text{Kr})}$	$\delta^\dagger$, %	P , Torr
4320	$9.75 \cdot 10^{-2}$	$2.77 \cdot 10^{-4}$	352.0	19.1	—	—	$1 \cdot 10^{-1}$
4326(1)	$4.69 \cdot 10^{-3}$	$1.30 \cdot 10^{-5}$	360.8	22.1	—	—	$1 \cdot 10^{-3}$
4326(2)	$5.51 \cdot 10^{-3}$	$1.60 \cdot 10^{-5}$	344.4	16.5	3.35	1.8	$1 \cdot 10^{-3}$
4413	$1.01 \cdot 10^{-2}$	$2.50 \cdot 10^{-5}$	404.0	36.7	—	—	$1 \cdot 10^{-3}$
4424	$8.88 \cdot 10^{-2}$	$2.72 \cdot 10^{-4}$	326.5	10.5	—	—	$1 \cdot 10^{-1}$
4427	$1.23 \cdot 10^{-2}$	$4.00 \cdot 10^{-5}$	307.5	4.1	—	—	$1 \cdot 10^{-1}$
4486	$8.73 \cdot 10^{-3}$	$2.50 \cdot 10^{-5}$	349.2	18.2	—	—	$1 \cdot 10^{-3}$
4511	$9.37 \cdot 10^{-2}$	$3.02 \cdot 10^{-4}$	310.3	5.0	3.13	-4.9	$1 \cdot 10^{-1}$
4519	$2.39 \cdot 10^{-3}$	$6.00 \cdot 10^{-6}$	398.3	34.8	2.99	-9.1	$1 \cdot 10^{-3}$
4615*	$1.01 \cdot 10^{-1}$	$3.48 \cdot 10^{-4}$	290.2	-1.8	—	—	$1 \cdot 10^{-1}$
4621**	$1.00 \cdot 10^{-1}$	$3.48 \cdot 10^{-4}$	287.4	-2.7	—	—	$1 \cdot 10^{-1}$

* gas sample taken without discharge in the reaction chamber;

** gas sample of the experiment, in which a discharge was directed beside the target;

† deviation of the isotope ratio from the atmospheric one.

time of the gas mixture under study. For both, the exact calibration of the scale of mass numbers of the mass-spectrometer and the determination of the amount of ^{40}Ar , we used the reference gas (a tracer) ^{38}Ar (the Ar-K method [see Ref. 78]). It was introduced in the measuring chamber of the mass-spectrometer in the amount of 1 ng, which ensured a signal of 10 V at the input resistance of a dc amplifier of 1×10^{11} Ohm.

In Table 8.1, we give results of the determination of the ratios of the basic isotopes of argon and krypton contained in the gas samples taken for studying and averaged over five counts. The isotope ratio for a chemical element in the analyzed sample can be judged by the ratio of their ion currents registered with a mass-spectrometer, since these currents are proportional to their content. In other words, the data in the fourth column of Table 8.1 can be considered to be the same as the ratios of the contents of isotopes ^{40}Ar and ^{36}Ar . Analyzing these data, we note that all the samples taken in experiments with “successful” damage of a target (i.e., the indication of an internal explosion action) are characterized by the enhanced isotope ratios (from 307.5 to 404.0) as compared to the ratio of their contents in the earth’s atmosphere (about 295.5 [Refs. 96–98]). At the same time, the ratio of the basic isotopes of argon in nonproductive experiments (Nos. 4615 and 4621) practically coincides with the atmospheric one. Values of the relative deviations of the registered isotope ratios of argon from their atmospheric ratio are given in the fifth column of Table 8.1.

We indicate the fact that the maximum anomalies of the isotope ratio of argon were registered in those cases where the pressure in the reaction chamber after the explosion of a target and the ion currents of its nuclides were minimum. On the contrary, these anomalies were minimum for a low rarefaction of the atmosphere of the reaction chamber and high ion currents. The mentioned correlations indicate, most likely, that a low rarefaction in the reaction chamber was conditioned by the enhanced content of atmospheric air in it and, hence, of atmospheric argon. Atmospheric argon caused high ion currents and significantly diluted the argon synthesized in nuclear reactions, anomalies of the isotope composition of which became inconspicuous due to a lower share in the total mixture with atmospheric argon. As for the reasons for the enhanced content of atmospheric air in the reaction chamber, we mention the following: a poor preliminary evacuation of the chamber, the entrance of air in it during puncturing a rubber seal by a sampler needle under sampling, and the degassing of surfaces of the body walls and details of the chamber in the cases where they were involved only partially to the process of high-voltage discharge.

As for the ratio of the basic isotopes of krypton, it was determined only for three gas samples (see Table 8.1). The largest deviation of their ratio from the atmospheric one equal to about 3.29 (Refs. 87–89) was registered in sample No. 4519, i.e., in the sample, in which we observe the maximum anomaly of the isotope ratio of argon, the minimum ion currents of the analyzed nuclides, and a high rarefaction in the reaction chamber. We also note the circumstance that the isotope ratio of krypton is shifted to the side of the minor nuclide as distinct from argon.

While analyzing the gas samples taken from the reaction chamber after the target explosion, we were interested in not only the isotope ratios of inert gases, but also their relative contents. In particular, we determined the contents of He, Ne, and Kr relative to nuclide ^{40}Ar (the basic nuclide of argon). As representatives of the studied inert gases, we chose their basic nuclides, namely ^4He , ^{20}Ne , and ^{84}Kr . The results of this investigation are given in Table 8.2. First of all, we consider the last row of Table 8.2. There, we present the shares of the amounts of basic nuclides of helium, neon, and krypton as compared to those of the basic nuclide of argon in air. They were calculated based on the data on the volume content of inert gases in air ($5.24 \times 10^{-4}\%$ He, $1.82 \times 10^{-3}\%$ Ne, $9.34 \times 10^{-1}\%$ Ar, and $1.14 \times 10^{-4}\%$ Kr [see Ref. 83]) and the abundance of their basic nuclides there (100% ^4He , 90.51% ^{20}Ne , 99.6% ^{40}Ar , and 57% ^{84}Kr [see Refs. 100, 101]). In experiments without damage of a target (No. 4621) and without discharge (No. 4615), the ion currents of all the nuclides of inert gases, except for ^{40}Ar , were practically at the background level (at most 3×10^{-5} nA). The data concerning samples

Table 8.2. Results of the analysis of the content of inert gases in the samples taken from the reaction chamber atmosphere.

Sample	$I(^4\text{He})$, nA	$\frac{I(^4\text{He})}{I(^{40}\text{Ar})}$	$I(^{20}\text{Ne})$, nA	$\frac{I(^{20}\text{Ne})}{I(^{40}\text{Ar})}$	$I(^{40}\text{Ar})$, nA	$I(^{84}\text{Kr})$, nA	$\frac{I(^{84}\text{Kr})}{I(^{40}\text{Ar})}$
4320	—	—	—	—	$9.75 \cdot 10^{-2}$	—	—
4326(1)	—	—	—	—	$4.69 \cdot 10^{-3}$	$2.00 \cdot 10^{-4}$	$4.26 \cdot 10^{-2}$
4326(2)	—	—	—	—	$5.51 \cdot 10^{-3}$	$2.00 \cdot 10^{-4}$	$3.63 \cdot 10^{-2}$
4413	—	—	—	—	$1.01 \cdot 10^{-2}$	—	—
4424	—	—	—	—	$8.88 \cdot 10^{-2}$	—	—
4427	$5.00 \cdot 10^{-5}$	$4.04 \cdot 10^{-3}$	—	—	$1.23 \cdot 10^{-2}$	$2.10 \cdot 10^{-4}$	$1.71 \cdot 10^{-2}$
4486	$5.40 \cdot 10^{-5}$	$6.19 \cdot 10^{-3}$	—	—	$8.73 \cdot 10^{-3}$	~ 0	~ 0
4511	$6.10 \cdot 10^{-5}$	$6.51 \cdot 10^{-4}$	$5.65 \cdot 10^{-4}$	$6.03 \cdot 10^{-3}$	$9.37 \cdot 10^{-2}$	$1.93 \cdot 10^{-4}$	$2.06 \cdot 10^{-3}$
4519	$9.20 \cdot 10^{-5}$	$3.85 \cdot 10^{-2}$	$1.59 \cdot 10^{-4}$	$6.65 \cdot 10^{-2}$	$2.39 \cdot 10^{-3}$	$2.45 \cdot 10^{-4}$	$8.45 \cdot 10^{-2}$
4615*	~ 0	~ 0	~ 0	~ 0	$1.01 \cdot 10^{-1}$	~ 0	~ 0
4621**	~ 0	~ 0	~ 0	~ 0	$1.00 \cdot 10^{-1}$	~ 0	~ 0

* gas sample taken without discharge in the reaction chamber;

** gas sample of the experiment, in which a discharge was directed beside the target;
in which a discharge was directed beside the target.

Nos. 4320, 4326(1), 4326(2), 4413, and 4424 should be considered qualitative, because we refined the procedure of determination of the relative content of helium, neon, and krypton using them. In all the remaining cases, the presented values are averaged over five counts derived by the extrapolation to the injection time of the studied sample into the measuring chamber of a mass-spectrometer.

Analyzing these data, we note the following. In the cases where the pressure in the reaction chamber after the explosion of a target was minimum, we registered that the gas samples contained the relative contents of the nuclides of the inert gases under study which exceeded considerably their relative contents in the atmosphere: by three orders of magnitude for ^{84}Kr [samples Nos. 4326(1), 4326(2), 4427, and 4519], by orders for ^4He (samples Nos. 4427, 4486, and 4519), and by 1 order for ^{20}Ne (sample No. 4519). As for the case of a medium rarefaction of the atmosphere in the reaction chamber (the dissolution of the synthesized inert gases by atmospheric air), we observed (sample No. 4511) only an insignificant excess of the relative

Table 8.3. Results of the determination of the isotope composition of krypton in the gas sample No. 4519.

Nuclide	Registered isotope ratio, %	Atmospheric isotope ratio*, %	Deviation of the isotope ratio, %
^{78}Kr	0.23	0.36	-36.1
^{80}Kr	2.58	2.28	13.2
^{82}Kr	11.52	11.58	-0.5
^{83}Kr	11.61	11.52	0.8
^{84}Kr	55.48	56.96	-2.6
^{86}Kr	18.58	17.3	7.4

* see Refs. 100, 101.

contents of helium and neon above the atmospheric ones. But it remained high for krypton as before (two orders). The discovered increase in the relative contents of helium, neon, and krypton can be explained only by their nucleosynthesis occurring upon the explosions of the targets.

In the gas sample No. 4519, we registered maximum anomalies of the isotope composition for argon (see Table 8.1) and the highest relative contents of helium, neon, and krypton (see Table 8.2). The indicated characteristics of this sample were derived at the minimum ion currents of the analyzed nuclides, and the very gas sample was taken from the reaction chamber under a high degree of rarefaction of its atmosphere. All these facts indicate that, in this sample, which is a mixture of the residual air with synthesized inert gases, the share of synthesized inert gases was maximum. In other words, sample No. 4519 turned out to be unique in some sense. This fact became the main reason for the execution of the isotope analysis of krypton in this sample.

The results of this analysis averaging over 10 counts are presented in Table 8.3. Analyzing these data, we indicate the considerable deviations of the contents of nuclides ^{78}Kr , ^{80}Kr , ^{84}Kr , and ^{86}Kr in the analyzed gas sample from their abundance in air. Moreover, the decreased contents were characteristic of nuclides ^{78}Kr and ^{84}Kr , whereas the contents of nuclides ^{80}Kr and ^{86}Kr were increased. As for nuclides ^{82}Kr and ^{83}Kr , their concentrations were close to their atmospheric abundance.

By studying the anomalies of the isotope composition of gaseous samples, we cannot but consider the question about the possible mechanisms of fractionation of isotopes. In the procedure of analysis of the isotope

composition of inert gases which was used by us, fractionation of isotopes can occur, in principle, in the following two stages: upon the evacuation of the reaction chamber and upon the sampling. By virtue of the fact that a lighter nuclide is more mobile, the evacuation of the reaction chamber can lead to the enrichment of its residual atmosphere by heavier nuclides, whereas the sampling can cause the enrichment of a gas sample by lighter nuclides. Both mechanisms bear a compensating character relative to each other. However, it is very difficult to judge, basing on the general consideration, the degree of their manifestation and how they compensate each other. To clarify this question, we appeal to experimental data. They indicate that the considered mechanisms of fractionation of isotopes do not manifest themselves. Indeed, in the case of sample No. 4615, we made no shot on a target and registered neither anomalies of the isotope composition of inert gases, nor their enhanced content relative to the atmospheric one, i.e., the effect of fractionation of nuclides was absent. Further, the anomaly of the isotope composition of argon in gas samples is shifted to the side of heavy nuclide ^{40}Ar , whereas the anomaly of the isotope composition of krypton is shifted, on the contrary, to the side of lighter nuclide ^{80}Kr . It is obvious that the working mechanism of fractionation of isotopes would lead to a shift of the isotope ratios of various inert gases to the same side. Finally, we note that the increased relative contents of helium, neon, and krypton as compared to the atmospheric ones were registered in sample No. 4519. However, if the efficient mechanisms of fractionation of nuclides would operate, then the relative contents of nuclides lighter and heavier than argon, would deviate to opposite sides.

Thus, we may conclude that the registered anomalies of the composition of gas samples taken from the reaction chamber after the explosion of a target are not a consequence of the fractionation of nuclides. We are sure that the deviations of the isotope compositions of argon and krypton and the enhanced relative contents of helium, neon, and krypton in the studied gas samples are a consequence of the running of nuclear reactions upon the explosions of targets in the hermetic reaction chamber.

8.1.2. Isotope Composition of Target Explosion Products

The gas component of target explosion products, for which the results of analysis of the isotope composition are presented in the above section, is an important, but insignificant part of the explosion products. Their most part is solid explosion products deposited in the form of thin layers on the surfaces of accumulating screens. However, as distinct to the situation with a gas component, the determination of the isotope composition of a solid part of explosion products is conjugated with a number of analytical difficulties.

These difficulties are caused, by the fact that solid target explosion products deposited on accumulating screens are inhomogeneous and multicomponent objects. Moreover, the amount of target explosion products captured by an accumulating screen is equal for now to several mg, and a regenerated target substance is only their part. All this creates some difficulties in the interpretation of mass-spectra (a lot of mass-peaks, their small amplitudes, and frequent case of their superpositions) registered from solid target explosion products with the purpose to determine the isotope composition of chemical elements contained in them.

At the same time, it should be noted that the mentioned difficulties are not insurmountable, though they present a number of complications upon the decoding of complex mass-spectra (see Refs. 99, 102–106). Indeed, the set of methods of mass-spectrometry contains sufficient number of standard tools to solve the problem of identification of overlapping mass-peaks which will be considered below upon the analysis of the results derived on their application. At this point, we indicate some aspects in the specificity of an attitude of the scientific community to results based on the analysis of complex mass-spectra. Experts in the field of mass-spectrometry believe that complex mass-spectra can be interpreted, and reliable results obtained on the basis of their analysis by people with sufficiently high qualifications. An analogous viewpoint is also supported by experts-experimenters using results of the method. However, they keep this viewpoint only if the results are referred to ordinary ones. In those cases where the results touch or contradict the fundamental physical ideas well-known in their field, experts-experimenters consider traditionally all such results as erroneous (arisen due to the incorrect identification of the mass-peaks of molecular ions) without any analysis of details of the decoding. The groundlessness of such an approach is obvious, but we should keep it in mind, to our regret.

In view of the formed situation, in order to avoid the obstruction of all the results of the determination of the isotope composition of solid target explosion products, we separate that part from them which is not connected, by its origin, with the analysis of complex mass-spectra registered on multi-component systems. We will consider the following two cases. First, we use pure chemical elements as the initial materials of targets and accumulating screens in explosive experiments and analyze the isotope composition of the basic chemical element of a material of the initial target in explosion products. Second, we determine the isotope composition of explosion products of a solid target by laser mass-spectrometry with a laser operating in the mode of modulation of the resonator quality. In the first case, we deal with the analysis of simple mass-spectra, in which the significant amplitudes are characteristic of only the mass-peaks related to the nuclides of one (or two)

base chemical elements of a material (materials) of the initial target and the accumulating screen depending on whether they coincide or differ. In the situation under study, the effect of superposition of mass-peaks from molecular and multicharged ions can be neglected by virtue of the fact that the total amount of all extrinsic chemical elements generating the mentioned mass-peaks is less by several orders than the content of the analyzed chemical element in the studied sample of nuclides. In the second case, the peaks of molecular ions are not observed in mass-spectra because a laser operating in the mode of modulated resonator quality induces the full dissociation of clusters due to a high temperature of the formed plasma (see Refs. 99, 105).

Isotope Composition of Base Chemical Elements of a Material of Targets in Explosion Products. In our investigation of the isotope composition of basic chemical elements of a material of targets in solid products of an explosion, we use a highly sensitive glow-discharge mass-spectrometer VG 9000 (VG Elemental, UK). As a holder, we used a tantalum cell for planar specimens without cooling which provided the diameter of the analyzed region to be 5 mm. The cathodic spraying of a specimen was realized in a glow discharge of Ar plasma with an energy of ions of 0.5 to 1.0 keV. The discharge current was 0.5 to 0.8 mA. The residual pressure in the chamber with specimens and in the spectrometer was at most 1×10^{-2} Pa and 1×10^{-5} Pa, respectively. For the ion beam, we used an accelerating voltage of 8 kV. The spectrometer possessed the range of analyzed masses 1 to 250, and its resolution in masses $M/\Delta M$ at the half-height of the Cu mass-peak was at the level of about 7000 to 9000. To register ion currents with the device, we used a photoelectric multiplier (the range of small currents) and a Faraday cup (the range of large currents). Their measurement time of the ion current of a nuclide was about 10 and 20 s, respectively.

All the studies considered in this section were carried out on accumulating screens with target explosion products deposited on them (see Fig. 8.1, b). Chemically pure copper served as a material for screens, and pure silver, iron, and lead were taken as the materials of initial targets. Screens had the form of disks of 0.5 to 1.0 mm in thickness and of 10 to 15 mm in diameter and served as a substrate. The studied layer of the material consisting of deposited target explosion products was on one side of their two surfaces (the working surface). The specimens were not subjected to any preparing procedures prior to the investigation and did not store for a long time. That is, they were taken as-received and were studied at once after the explosion-based experiments.

During the study, the 5-mm analyzed spot of a mass-spectrometer was located directly on the surface of the layer of target explosion products.

In this case, the direction of analysis was normal to the working surface of an accumulating screen. The mentioned scheme of analysis was chosen with the purpose of deriving the maximum signal upon the registration of the ion currents of the nuclides of studied chemical elements and, hence, of enhancing the accuracy of the determination of their isotope composition. As noted above, this scheme is very useful for the determination of the isotope composition of the basic chemical element of the initial target material due to the simple interpretation of mass-spectra. Due to the intense signal of ion currents, the scheme of analysis can be also used upon the determination of the isotope composition of minor chemical elements in target explosion products. However, its application in the mentioned case is limited by the following circumstance. Because of an insignificant thickness of a layer of explosion products, the scheme did not allow us to carry out the analysis of the isotope composition of chemical elements in a sufficiently wide mass range. On the etching rate of about $0.1\text{ }\mu\text{m}/\text{min}$ and the thickness of a layer of explosion products of 3 to $4\text{ }\mu\text{m}$, we were able to correctly perform the isotope analysis only for some chemical elements.

As for the use of this scheme for the analysis of the isotope composition of the basic chemical element of a target material, it is connected with overcoming a number of difficulties, though it has some advantages. In its proper application, the main problem is the inhomogeneity of the composition of a layer of explosion products. By virtue of that the registration of ion currents of nuclides with a glow-discharge mass-spectrometer VG 9000 runs in a successive mode (from a low-mass isotope to a high-mass isotope), any inhomogeneity of the etched layer can induce a distortion of the isotope ratios of an analyzed chemical element. Relative to the studied object, we may distinguish two types of inhomogeneities: an inhomogeneity in the plane parallel to the surface of the investigated layer and that inside a specimen. Due to the fact that the diameter of the analyzed region considerably exceeded the size of a characteristic structural element of explosion products, the first-kind inhomogeneity can be well averaged.

The second-kind inhomogeneity was caused by a change of the structure of the investigated layer in depth. First of all, large gradients of the layer composition were observed on its external surface and the boundary with the accumulating screen. Moreover, narrow regions with composition gradients were also located in the very bulk of a layer of explosion products. Their origin is related to the circumstance that the explosion products contain not only the regenerated, but also initial target substance. A spatial separation of the mentioned components under the dispersion of the target substance after the explosion was defined by its geometry resulted in their layerwise position on the accumulating screen. The boundary

between them was indicated by an internal region with variable composition. In other words, the target explosion products deposited on accumulating screens possessed a layered inhomogeneity.

We can deal with the last problem only in a single way: moving in the depth of a specimen during the etching, we should choose its almost homogeneous sections. We accomplished this in the following manner. We set a subsequent collection of windows, on the scanning of which a mass-spectrometer registered the ion currents of the nuclides of chemical elements under study. In this collection, we include, besides the windows for the nuclides of a single analyzed chemical element, additional windows for the nuclides of a number of other reference chemical elements, whose ion currents were used with the purpose of a control over the homogeneity of the etched section of a layer of explosion products and the stability of working parameters of the setup. Under a fixed position of the analyzed area on the accumulating screen, we performed several successive continuous scanning runs on the chosen collection of windows until the bottom of the etching crater reaches the substrate (the accumulating screen). Usually, we performed four such scans and thoroughly timed the whole measurement procedure. The control was realized by the comparison of the ion currents registered in the successive cycles of scanning for each nuclide.

To control the stability of the glow discharge combustion, we used the ion currents of nuclides ^{36}Ar (Ar plasma) and ^{181}Ta (the tantalum cell of the holder of specimens). Estimating the changes in the ion currents of nuclides ^1H , ^{12}C , ^{14}N , and ^{16}O which belong to the chemical elements entering the typical contaminations, we can judge the etching of the surface layer and the reaching of the substrate by the bottom of the etching crater. The control over the homogeneity of the bulk of the investigated layer was realized by using the ion currents from the nuclides of the basic chemical element of a target material and an accumulating screen and of several chemical elements contained in significant amounts in explosion products. A section of the bulk of a layer was considered homogeneous by its composition, if the ion currents of reference nuclides from it differed in two successive cycles of scanning by at most 5%. Sections not satisfying the indicated condition were excluded from the consideration.

Some words should be said about the chosen criterion of homogeneity of a section of the bulk of a layer of target explosion products. In the used scheme of analysis, the duration of one full cycle of scanning usually exceeded the duration of registration of ion currents of the collection of all nuclides of the basic chemical element of a target material in the cycle by more than one order. Hence, by assuming that the composition varies linearly in this section, we find that it has changed by at most 0.5% for the time of

registration of ion currents of the nuclides of an analyzed chemical element. It is a small value upon the analysis of the isotope composition of a chemical element. Moreover, this inhomogeneity was corrected to a significant extent. First, by the timing of the measurement procedure and by a change of the ion currents of reference nuclides in two successive cycles of scanning, we determined the rates of their variations. Then with regard to this rate and the timing, we corrected the ion currents of nuclides of the basic chemical element of a target material which were registered in this section.

Finally, we indicate one more circumstance, by which we were guided upon the selection of data for the calculation of the isotope ratios for the basic chemical element of a target material. As noted above, we used two detectors for the registration of ion currents by a mass-spectrometer: a photoelectronic multiplier and a Faraday cup which allowed us to cover, respectively, the range of small and large currents. In those cases where a significant difference in the abundances of the nuclides of an analyzed chemical element (targets made of iron or lead) took place, their ion currents could fall in the working range of different detectors. In this situation upon a low accuracy of the consistency of detectors, some distortions reveal themselves in measurements of the ion currents of nuclides. These distortions affect, in turn, the results of the determination of its isotope composition. With the purpose of avoiding permanent verifications and adjustments of the accuracy of the consistency of detectors or eliminating a potential source of errors in the determination of the isotope composition, we analyzed only those measurements, in which the ion currents of all nuclides of the analyzed chemical element were registered by one of the detectors.

To determine the isotope composition of silver in explosion products of silver targets, we studied five accumulating screens. The measurements were carried out at the center and on the periphery of each accumulating screen. Due to the closeness of values of the concentration of the silver nuclides, their ion currents registered on only one of the detectors in all cases of the performed measurements. Homogeneous sections of the bulk of a layer of explosion products which were suitable for the analysis of the isotope composition were separated only at the center and on the edge of accumulating screen No. 109 and at the center of accumulating screen No. 55. The results of the determination of the isotope composition of silver in the indicated measurements are given in Table 8.4.

Considering the data presented in Table 8.4, we can note, first of all, that the significant deviations of the isotope composition of silver from the natural abundance of its isotopes were registered only on specimen No. 109. Moreover, the mentioned anomalies occurred on it in both measurements: at the center and on the edge of the accumulating screen. As for their behavior,

Table 8.4. Results of the determination of the isotope composition of silver on specimens No. 55 and 109 in the experiments with silver targets and copper accumulating screens.

No. of a speci- men	Nuclide	1st scanning		2nd scanning		3rd scanning		Natural abun- dance*, %
		Conc., %	δ , %	Conc., %	δ , %	Conc., %	δ , %	
55 (center)	^{107}Ag	52.25	0.79	52.02	0.35	52.08	0.46	51.839
	^{109}Ag	47.75	-0.85	47.98	-0.38	47.92	-0.50	48.161
109 (center)	^{107}Ag	53.80	3.78	52.98	2.20	51.52	-0.62	51.839
	^{109}Ag	46.20	-4.07	47.02	-2.37	48.48	0.66	48.161
109 (edge)	^{107}Ag	55.19	6.46	55.69	7.43	52.17	0.64	51.839
	^{109}Ag	44.81	-6.96	44.31	-8.00	47.83	-0.69	48.161

* see Refs. 100, 101.

we can see some regularities. In particular, a special attention must be given to the fact that the anomalies of the isotope composition of silver are localized in both cases in the upper region of a layer of explosion products (see the 1st and 2nd cycles of scanning, Table 8.4). With moving in depth and approaching the layer bottom, they disappear (see the 3rd cycle of scanning, Table 8.4), and the isotope composition of silver tends to the natural one. We also indicate the circumstance that, in all registered cases of anomalies, the isotope composition of silver is shifted to the same side: to its nuclide with a lesser mass, i.e., with a lesser content of neutrons. In this connection, we noted also that small anomalies on specimen No. 55 have the same character. This circumstance allows us to assume that they are, more probably, low but significant deviations rather than the errors of measurements. As for the level of the registered anomalies, we may note two points: the range of its variation is in the interval from 2% to 8%, and its values are considerably higher on the edge of the accumulating screen, than at its center.

We try to comment on some indicated regularities. We start from the localization of anomalies of the isotope composition of the basic chemical element of a target material in the upper region of a layer of explosion products deposited on the accumulating screen. It points out unambiguously to that the upper and lower parts of the bulk of explosion products should be associated to a significant extent with a regenerated substance of the target and the initial target substance, respectively. Such a character of the location of these components of explosion products testifies to that they were spatially separated upon dispersion of the target substance, i.e., the regenerated target substance was inside a spherical layer of the initial target material. It is obvious that the dispersion geometry of the target substance

must inherit a geometry of the origination of the very explosion. In other words, we can conclude that the explosion origination zone is located in the inner region of the head part of a target, in which the energy of a beam of subrelativistic electrons from the vacuum diode is concentrated in the process of high-voltage discharge, and it is surrounded by the shell of the initial target material.

As for the stable regularity of deviations of the isotope composition of silver in the regenerated target substance to the side of its nuclide with a lesser content of neutrons, it is difficult to uniquely interpret it. From the purely logical viewpoint, we may consider two hypothetical versions of the development of events. On the one hand, we can come to such a result if we assume only the partial regeneration of an initial target substance occurs in the explosion origination zone. In such a situation, it is natural to assume the different stability or inclination of both nuclides of silver to nuclear transformations. In other words, there are some reasons to expect that the degrees of the “burning” of the isotopes of silver in running nuclear reactions will be different. It is obvious that this could lead to a difference of isotope compositions of that part of silver from the explosion origination zone, which underwent a transmutation, and the remaining part which was not transmuted. The indicated difference must bear a complementary character. Hence, if we register a shift of the isotope composition to the side of nuclide ^{107}Ag in the fraction of silver not undergone a nuclear transformation, its “burned” fraction must be enriched by nuclide ^{109}Ag .

On the other hand, we can also arrive at this result if we make assumption about the full regeneration of the target substance contained in the explosion origination zone. However, in this case, we must consider all silver contained in the transmuted target substance to be newly generated with the ratio of nuclides registered in experiments. Then, in turn, we should recognize the following fact which seemed rather fantastic: in the composition of a newly generated substance in the explosion origination zone, the dominant component is, though with other isotope composition, the same basic chemical element of the initial target material. In other words, the substance regenerated in the explosion remembers, as if, its prehistory and inherits, to a considerable extent, the parent signs. It seems to us intuitively that the evolution of the target substance in the explosion follows just this scenario.

Finally, we present some reasoning about the observation that the great anomalies of the isotope composition of the basic chemical element of a target material are registered not at the center, but on the periphery of an accumulating screen. In this connection, we note that the density of the flow of a hot substance transferred by explosion products decreases

proportionally to the squared distance from the explosion origination zone. Then it follows from the geometry of the relative position of the head part of a target and an accumulating screen, that the density of the mentioned flow at the center of the screen exceeds its density on the periphery of the screen by more than one order of magnitude. This circumstance can yield that, after the deposition of the first portion of explosion products (mainly the initial target substance) onto the central region of the screen and its heating, the condensation of subsequent portions (mainly a regenerated target substance) will be simultaneously accompanied by their removal through the evaporation and mixing with the earlier deposited lower layers. The mentioned lesser anomalies of the isotope composition at the central region of the accumulating screen could be a consequence of the described processes.

The determination of the isotope composition of iron in the explosion products of iron targets was carried out on four accumulating screens. Like in the case of the experiments with silver targets, we studied both, central and peripheral regions on each screen. According to the criterion of homogeneity of sections of the bulk of a layer of explosion products, we chose only three from the totality of measurements as suitable for the analysis. However, in two of them, the ion currents of the basic nuclide of iron and its rare nuclides were registered by different detectors. Therefore, they were also excluded from the consideration. Thus, as a result of the selection, we retained only one measurement which was carried out on the edge of accumulating screen No. 116. The results of the determination of the isotope composition of iron for this measurement are given in Table 8.5.

It is seen from the data presented in Table 8.5 that the isotope composition of iron transferred onto an accumulating screen as a result of the explosion of a target is considerably different from that characteristic of the

Table 8.5. Results of the determination of the isotope composition of iron on specimen No. 116 in the experiments with iron targets and copper accumulating screens.

No. of a speci- men	Nuclide	1st scanning		2nd scanning		3rd scanning		Natural abun- dance*, %
		Conc., %	δ , %	Conc., %	δ , %	Conc., %	δ , %	
116 (edge)	^{54}Fe	5.42	-6.55	5.81	0.17	5.73	-1.21	5.80
	^{56}Fe	91.71	-0.01	91.22	-0.55	91.31	-0.45	91.72
	^{57}Fe	2.51	14.09	2.59	17.73	2.60	18.18	2.20
	^{58}Fe	0.36	28.57	0.38	34.29	0.36	28.57	0.28

* see Refs. 100, 101.

natural abundance of its nuclides. First of all, this manifests in a significant increase in the content of nuclides ^{57}Fe and ^{58}Fe in the explosion products and in some decrease in the content of nuclide ^{54}Fe . The last is referred, in particular, to the upper part of the bulk of explosion products (see the 1st cycle of scanning). As for the basic nuclide ^{56}Fe , its content in explosion products is quite close to its natural abundance. Thus, the isotope composition of the regenerated iron is shifted, as distinct from the situation with silver, to the side of its nuclides with high content of neutrons. We also note that these deviations bear a regular character during all three cycles of scanning. In this case, their values are in the interval from 14% to 34%.

To study the isotope composition of lead in the explosion products of lead targets, we used 10 accumulating screens. These studies were performed earlier than those concerning the determination of the isotope compositions of silver and iron in experiments with silver and iron targets. On these specimens, we developed and adjusted the method of determination of the isotope composition of the basic chemical element of a target material in explosion products deposited on an accumulating screen. The problems we succeeded to solve in experiments with silver targets did not appear immediately. First, we posed and solved simpler problems: in any region of an accumulating screen (at the center or on the periphery), to find at least one homogeneous section of the bulk of explosion products and to register the anomalies of the isotope composition of lead there. Moreover, from the analytical viewpoint, lead turned out to be a complex object as for the determination of the isotope composition. We get that the ion currents of its rare nuclide ^{204}Pb and the rest of nuclides (^{206}Pb , ^{207}Pb , and ^{208}Pb) were almost always registered by different detectors. For example, from many measurements performed at the center and on the edge of accumulating screens, we separated 11 homogeneous sections of the bulk of a layer of explosion products suitable for the analysis of the isotope composition. But we succeeded only twice in registering the ion currents of all the nuclides of lead by one detector (a Faraday cup). The results of these measurements performed in the central region of accumulating screens No. 103 and 88 are presented in Table 8.6.

Prior to the analysis of the data given in Table 8.6, we note the following. We compared the registered isotope composition of lead in explosion products not with its natural abundance (see the last column in Table 8.6), but with the isotope composition of the initial target material (see the penultimate column in Table 8.6). This circumstance is induced by the fact that the isotope composition of lead not undergone any treatment can vary considerably, in the general case, from that given in reference books as natural. The data on the isotope composition of the initial substance of lead

Table 8.6. Results of the determination of the isotope composition of lead in the explosion products on specimens No. 88 and 103 in the experiments with lead targets and copper accumulating screens.

No. of a specimen	Nuclide	1st scanning		Conc. in the standard, %	Natural abund.*, %
		Conc., %	δ , %		
103 (center)	^{204}Pb	1.38	2.96	1.34	1.4
	^{206}Pb	25.17	2.76	24.49	24.1
	^{207}Pb	21.40	-0.11	21.42	22.1
	^{208}Pb	52.06	-1.33	52.76	52.4
88 (center)	^{204}Pb	2.10	56.53	1.34	1.4
	^{206}Pb	35.50	44.94	24.49	24.1
	^{207}Pb	17.41	-18.73	21.42	22.1
	^{208}Pb	45.00	-14.71	52.76	52.4

* see Refs. 100, 101.

targets, which are given in Table 8.6, were derived on a device VG 9000. They are averaged over the results of 4 measurements.

As for the results of the determination of the isotope composition of lead in the explosion products of lead targets, we note the following. For the measurement performed at the center of accumulating screen No. 103, the isotope composition of lead is rather close to that of the initial substance of targets. As distinct from the previous case, the measurement at the center of accumulating screen No. 88 revealed a giant anomaly of the isotope composition of lead. Moreover, the anomalous content is characteristic of all nuclides of lead: the level of deviations is about 45% to 56% for light nuclides and 15% to 19% for heavy ones. On the whole, the anomaly can be characterized as a strong shift of the isotope composition to the side of nuclides with a lesser content of neutrons.

We now return to those measurements on a selected homogeneous section of the bulk of a layer of explosion products suitable for the analysis of the isotope composition which were set aside, because the ion current of rare nuclide ^{204}Pb was registered by the other detector. As mentioned above, the number of such measurements was sufficiently high. To avoid the loss of such great amount of data, we reserved them for the rest nuclides of lead: ^{206}Pb , ^{207}Pb and ^{208}Pb . In this connection, we meet the question

about the form of their representation. In similar cases, the data are usually given as the ratios of the measured contents of nuclides to the content of the basic nuclide. However, we consider such form of a representation of results not quite convenient. The point to note is that we will have the data only on two nuclides instead of on three, since the mentioned ratio for the basic nuclide will be always equal to one. In our opinion, it is more informative to represent these data so that the sum of contents of three considered nuclides of lead (^{206}Pb , ^{207}Pb , and ^{208}Pb) be 100%. That is, we should normalize the content of each nuclide on the sum of contents of three indicated nuclides, rather than on the content of a basic nuclide. The results of nine remaining measurements in this form are presented in Table 8.7.

Commenting upon the results given in Table 8.7, we note that the isotope composition of lead close to that in the initial substance of targets (see specimen No. 104, edge) was registered only in one case from all measurements. In eight remaining measurements, we have the anomalous isotope composition of lead. The maximum level of these anomalies reaches 8% to 9%. We indicate that the last numbers demonstrate sufficiently large deviations in the composition of nuclides, whose contents are about several tens of percents. It is remarkable that, in all these cases except one (specimen No. 104, center), the isotope composition of lead shifts, as before, to the side of its nuclides with lesser content of neutrons.

Finally, we mention once again that the anomalies of the isotope composition registered here cannot be a consequence of an erroneous interpretation of mass-spectra because of the interference of the mass-peaks of nuclides of the analyzed chemical element with the mass-peaks of molecular complexes. This is related to the circumstance that the content of the analyzed element in the studied objects was higher by several orders than the content of chemical elements in them, which would form such complexes and would considerably influence the intensity of the mass-peaks of nuclides of the analyzed element. In particular, this is referred to anomalies of the contents of widespread nuclides (^{107}Ag , ^{109}Ag , ^{206}Pb , ^{207}Pb , and ^{208}Pb) in the studied objects. Thus, the discovered anomalies of the isotope composition of a number of chemical elements indicate that the target explosion products deposited on accumulating screens contain a substance of the artificial origin.

Investigation of the Isotope Composition of Solid Products of Explosions by Laser Mass-Spectrometry. As in the previous case, we took a layer of solid target explosion products deposited on accumulating screens of a standard size (see Fig. 8.1, b) as the object of studies upon the analysis of an isotope composition by laser mass-spectrometry. As the initial

Table 8.7. Results of the determination of the isotope composition of lead in explosion products in the experiments with lead targets and copper accumulating screens (the data are normalized on the sum of contents of the analyzed nuclides).

No. of a speci- men	Nuclide	1st scanning		2nd scanning		Conc. in the standard, %	Natural abund.*, %
		Conc., %	δ , %	Conc., %	δ , %		
104 (edge)	^{206}Pb	24.54	-1.12			24.82	24.44
	^{207}Pb	21.53	-0.81			21.71	22.41
	^{208}Pb	53.93	0.83			53.48	53.14
103 (edge)	^{206}Pb	26.83	8.11			24.82	24.44
	^{207}Pb	22.00	1.35			21.71	22.41
	^{208}Pb	51.16	-4.33			53.48	53.14
102 (edge)	^{206}Pb	25.91	4.41	26.00	4.75	24.82	24.44
	^{207}Pb	22.26	2.53	22.42	3.26	21.71	22.41
	^{208}Pb	51.83	-3.09	51.58	-3.55	53.48	53.14
90 (center)	^{206}Pb	27.09	9.13			24.82	24.44
	^{207}Pb	21.88	0.80			21.71	22.41
	^{208}Pb	51.03	-4.58			53.48	53.14
84 (center)	^{206}Pb	26.10	5.15	25.62	3.22	24.82	24.44
	^{207}Pb	22.15	2.04	21.80	0.43	21.71	22.41
	^{208}Pb	51.75	-3.24	52.58	-1.69	53.48	53.14
104 (center)	^{206}Pb	24.27	-2.23	22.50	-9.33	24.82	24.44
	^{207}Pb	20.23	-6.81	20.96	-3.45	21.71	22.41
	^{208}Pb	55.50	3.78	56.53	5.71	53.48	53.14

* see Refs. 100, 101.

substance of screens, we usually used chemically pure copper and rarely used Ag, Ta, Au, Pt, etc. Targets were made of light, medium, and heavy, chemically pure metals with atomic masses in the range from 9 to 209. All specimens were studied at once after the explosion-based experiments and in the state, in which they were derived in these experiments.

The determination of the isotope composition of target explosion products deposited on the working surface of accumulating screens was realized on a nonserial laser mass-spectrometer designed and manufactured at Taras Shevchenko Kiev National University. In this mass-spectrometer, the laser operating in the mode of modulation of the resonator quality serves as a means of the microsampling from a specimen. The operation of the laser in this mode is characterized by the fact that the emitted energy is somewhat less than that in the mode of free generation. However, in the former case at the expense of a considerable reduction in the duration of a laser pulse ($\tau \sim 10^{-8}$ s), the density of a released power increases approximately by 10^3 times ($q \sim 10^8$ to 10^{10} W/cm², the heating rate is $\sim 10^{10}$ K/s) (see Refs. 99, 105). As a result, the amount of the evaporated mass is small ($\sim 10^{-10}$ g). In this case, the dominant share of the laser pulse energy (94% to 96%) is not spent on the heating of a specimen substance, but is transferred to the removed products. This transits them in the state of fully dissociated (the absence of the peaks of clusters in mass-spectra) and practically fully ionized plasma ($T \sim 10^4$ K) (see Ref. 105).

The analysis of masses in a microsample evaporated by a laser ray was carried out with a time-of-flight mass-spectrometer (see Ref. 99). The principle of separation of masses in this device is based on that the time of flight of ions τ in the vacuum drift space along a fixed distance L from the place of their origination to a collector depends on the ratio of the mass of an analyzed ion M to its charge q and is defined by the formula:

$$\tau = L(M/2qU_0)^{1/2}, \quad (8.1)$$

where U_0 is the accelerating voltage (see Ref. 105). Hence, during the movement from the source to the collector, ions are separated in bunches, respectively, to the ratio of their mass to charge M/q . The first to arrive at the collector are the bunches of light ions, and then the heavier ones approach it, successively. There, they are registered by a detector and are displayed on the monitor of an oscilloscope in the form of peaks at the positions corresponding to the ratio M/q .

The device allows us to analyze masses in the range from 1 to 300 a.u.m. and possesses the sensitivity of $10^{-4} \times 10^{-5}\%$ and spectral resolution of $M/\Delta M \sim 600$. The locality of such a method is defined by the geometric size of a microsample taken from a specimen. In the used device, microsamples have form of a disk of 100 to 150 μm in diameter and 1.0 to 1.5 μm in thickness.

Upon the study of the isotope composition of explosion products, the sampling was carried out directly from the layer deposited on the working surface of an accumulating screen. That is, the scheme of analysis was the

same as in the case of the determination of the isotope composition in explosion products of the basic chemical element of a target material. As was mentioned above, this scheme of analysis ensured an intense signal upon the registration of the ion currents of nuclides of the studied chemical elements and promoted an enhancement of the accuracy of the determination of their isotope composition. The indicated circumstance was important in the considered case of the analysis of the isotope composition of minor chemical elements because of their low content in target explosion products.

By virtue of the fact that the studied specimens had axial symmetry, we assumed that the chemical and isotope compositions of a layer of explosion products is independent of the choice of a radial direction onto the accumulating screen (of the azimuth angle ϕ) and is only a function of the distance from the center of a specimen to a selected place of analysis (r). Therefore, upon the investigation of specimens, we usually chose several most typical radial directions and carried out the sampling for the analysis at six to eight points located more or less uniformly along each direction. We also undertook the attempts to determine the character of variations in the isotope compositions of target explosion products by moving from the surface of a specimen by means of repeated samplings for the fixed place of analysis. However, usually upon the repeated measurements, a considerable amount of a substrate material enters into a microsample, which testifies to that the thickness of the layer of products of the reaction is at most 2 to 3 μm .

We will also discuss the influence of chemical inhomogeneities of the investigated layer on the results of the determination of its isotope composition. First of all, we note that the inhomogeneity in a plane parallel to the investigated layer surface was well-averaged because the diameter of the analyzed region significantly exceeded the size of a characteristic structural element of explosion products. As for the lamellar inhomogeneity, it can lead to a distortion of results of the determination of the isotope ratios of an analyzed chemical element only if the mode of successive sampling is in use, i.e., if the analyzed nuclides referred to the same chemical element are taken for the registration of their ion currents from sections of the analyzed layer which are positioned at different depths. In the method of laser mass-spectrometry in the successive mode, there occurs only the registration of ion currents (a time-of-flight analyzer of masses). However, the analyzed nuclides of all chemical elements are taken from the same volume (see Refs. 99, 105). This yields that the results of measurements in the analysis scheme under consideration have sense of the mean isotope composition of a selected microsample. The last factor means that if we registered an anomaly of the isotope composition of any chemical element, the anomaly cannot be

a consequence of the concentration gradient of this chemical element with the natural abundance of its nuclides.

Upon the analysis of the isotope composition, we first decoded a mass-spectrum registered on the investigated specimen. Then by selecting the time parameters of the scanner, we separated a fragment of the spectrum on the monitor of an oscilloscope with peaks concerning the isotopes of an analyzed chemical element. As a calibration curve, we supplied an analogous fragment of the mass-spectrum of the same chemical element with the natural distribution of isotopes, which was registered on the standard, to the upper part of the display and made comparison. Since we used pure chemical elements or their simple compounds as the substance of standards, the amplitude of reference mass-peaks frequently considerably exceeded the amplitudes of the appropriate mass-peaks of an analyzed chemical element in a specimen, whose content was usually at a significantly lower level. Therefore, for the convenience of the comparison, we decreased the amplitudes of the reference mass-peaks on the display of the oscilloscope by 10^1 to 10^2 times relative to the amplitudes of the studied mass-peaks.

In the discussion of the question of the comparison of the corresponding fragments of the mass-spectra of a standard and a specimen under study, we indicate two more circumstances. In the cases where the content of an analyzed chemical element in a specimen is low, a noticeable modulated structure (a splitting of peaks) can sometimes appear in mass-peaks of its nuclides, and the position of the very mass-peaks in the spectrum can undergo a certain shift relative to the calibration curve. In this case, the mentioned phenomena behave themselves so that, with increase in the content of an analyzed chemical element in a specimen, the frequency grows, and the modulation amplitude of mass-peaks of its nuclides decreases. They become smoother, and their shift also decreases, that is, the manifestations of both effects become less pronounced and then disappear. The mentioned effects are defined by the ion–electron and ion–ion interactions in the plasma of the region where the ion bunches of a laser time-of-flight mass-spectrometer are separated. Their nature and mechanisms are described in Refs. 107, 108.

It is necessary to note that the splitting of mass-peaks can lead to a significant decrease of their amplitudes due to their high-frequency modulation. However, it should indicate that the area under modulated mass-peaks preserves in this case, because the total charge of the ion bunch forming the analyzed mass-peak does not change upon its interaction with electrons of the plasma and other ion bunches in the region where ion bunches of a spectrometer are separated. With regard to the above-written, we used the area under mass-peaks instead of their amplitudes in the determination of

the isotope composition of target explosion products in those cases where mass-peaks were split.

Studying the target explosion products deposited on the surface of accumulating screens, we registered the deviations of the isotope composition from the natural abundance of nuclides for such chemical elements as S, Cl, Ca, Cr, Fe, Ni, Cu, and others. For various specimens, the characters of the discovered deviations of the isotope composition for a certain chemical element were different. However, in the first approximation in the scope of one specimen, the isotope composition can be considered to be invariable. In other words, the anomalous isotope composition of an analyzed chemical element was approximately the same independently of the place of registration of the mass-spectrum on the surface of the investigated specimen, though we found the regions with maximum content of a chemical element possessing the anomalous isotope composition in the studied specimens. The mentioned enriched regions were usually located at a distance from the center of a specimen which was equal approximately to its half-radius, i.e., on the edge of a crater (see Fig. 8.1, b).

The above-described character of the distribution of an analyzed chemical element over the surface of a specimen testifies to that the registered deviations of its isotope composition from the natural abundance of its nuclides are not a consequence of the separation of isotopes of the natural origin. For the effect of separation of isotopes, we would observe the clear dependence of the value of a deviation of the isotope composition on the distance between a place of analysis on the surface and the specimen center. We add that it is difficult to invent mechanisms which would provide efficient separation of the isotopes of a chemical element under a spherically symmetric explosion.

Below, we will illustrate the brightest cases of anomalies of the isotope composition of chemical elements which were discovered by laser mass-spectrometry upon the study of target explosion products. In the lower part of Fig. 8.2, a, we show a fragment of the mass-spectrum which was registered on the explosion products of a lead target deposited on copper accumulated screen No. 3414 and contained the collection of mass-peaks of the isotopes of Cl. The upper part of this figure demonstrates the relevant fragment of the mass-spectrum (it is reduced relative to the lower fragment by 50 times by vertical) of Cl registered on the KCl standard with natural abundance of its nuclides (see also Table 8.8). In the standard mass-spectrum, we see two peaks referred to nuclides ^{35}Cl and ^{37}Cl with the approximate ratio of the amplitudes of 3 : 1, respectively. In the spectrum registered on the specimen, we must see the mass-peak of ^{37}Cl for the indicated ratio of the amplitudes of mass-peaks of the isotopes of Cl. However, it is absent in the presented

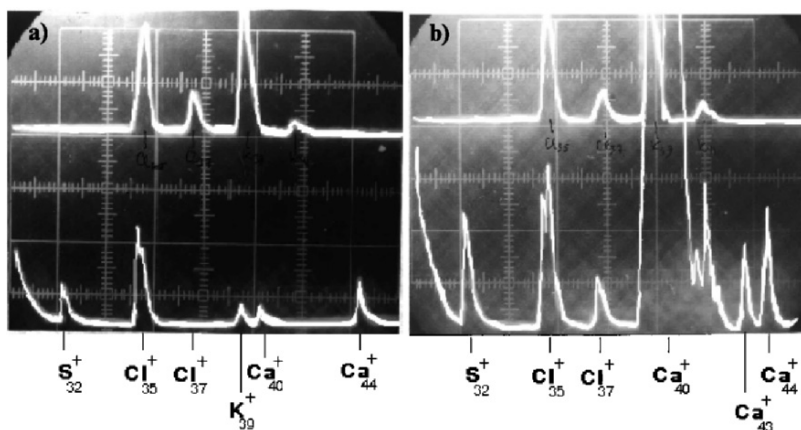


Fig. 8.2. Comparison of the isotope compositions of Cl (a) and Ca (a, b) contained in explosion products with their natural isotope composition.

fragment of the spectrum. This fact means that the chemical element Cl is represented here only by one nuclide, ^{35}Cl .

In this mass-spectrum, we can also observe the violation of the isotope composition of Ca. Natural calcium has six stable nuclides (see Table 8.8). Its main nuclide is ^{40}Ca . In the natural calcium, the content of nuclide ^{44}Ca is approximately 50 times lower than that of the basic nuclide ^{40}Ca , whereas the concentration of other nuclides of Ca is yet less. However, the mass-spectrum registered on explosion products of a lead target and given in Fig. 8.2, a, contains the mass-peak of nuclide ^{44}Ca , whose amplitude exceeds the amplitude of the mass-peak of the basic nuclide ^{40}Ca by more than double.

The fragment of the mass-spectrum presented in the lower part of Fig. 8.2, b was registered on the explosion products of a copper target on copper accumulating screen No. 3417. The corresponding standard mass-spectrum, as in the previous case, is reduced by 50 times by vertical. In the given mass-spectrum, we again see a violation of the isotope composition of Ca. Here, the ratio of amplitudes of the mass-peaks of nuclides ^{40}Ca and ^{44}Ca is close to the natural one, whereas the amplitude of the peak of nuclide ^{43}Ca is considerably higher.

According to the natural ratio (see Table 8.8), the amplitude of the mass-peak of nuclide ^{43}Ca must be approximately 15 times less than that of the mass-peak of nuclide ^{44}Ca . But they are comparable on the registered mass-spectrum.

In the upper part of Fig. 8.3, a, we give a fragment of the mass-spectrum with natural distribution of the isotopes of sulfur (it is reduced

Table 8.8. Data on the natural abundance of nuclides of some chemical elements.

Nuclide	Mass number, a.u.m.	Abundance (Refs. 100, 101), %	Nuclide	Mass number, a.u.m.	Abundance (Refs. 100 101), %
^{32}S	31.972070	95.020	^{54}Fe	53.939610	5.800
^{33}S	32.971460	0.750	^{56}Fe	55.934940	91.720
^{34}S	33.967870	4.210	^{57}Fe	56.935390	2.200
^{36}S	35.967080	0.020	^{58}Fe	57.933280	0.280
^{35}Cl	34.968850	75.770	^{58}Ni	57.935350	68.270
^{37}Cl	36.965900	24.230	^{60}Ni	59.930790	26.100
^{40}Ca	39.962590	96.941	^{61}Ni	60.931060	1.130
^{42}Ca	41.958620	0.647	^{62}Ni	61.928350	3.590
^{43}Ca	42.958770	0.135	^{64}Ni	63.927970	0.910
^{44}Ca	43.955490	2.086	^{63}Cu	62.929600	69.170
^{46}Ca	45.953690	0.004	^{65}Cu	64.927800	30.830
^{48}Ca	47.952530	0.187			

by 10 times relative to the lower fragment). The lower part shows the corresponding fragment of the mass-spectrum registered on specimen No. 3414 (a lead target and a copper accumulating screen). In the Nature, sulfur has four isotopes (see Table 8.8). ^{32}S is its basic nuclide. The content of nuclide ^{34}S in natural sulfur is more than 20 times less than that of its basic nuclide, and the concentration of its other isotopes is yet less. In the mass-spectrum registered on explosion products, together with the mass-peak of its basic nuclide, we see the mass-peak of nuclide ^{33}S . But the mass-peak of nuclide ^{34}S is absent, though its amplitude would exceed the amplitude of the mass-peak of nuclide ^{33}S by at least five times according to the natural abundance.

Continuing on we refer to Fig. 8.3, b (a platinum target and copper accumulating screen No. 3013), Fig. 8.4, a (a gold target and copper accumulating screen No. 2783), and Fig. 8.4, b (a copper target and copper accumulating screen No. 2569) show the anomalies of the isotope composition of Ni. In all the cases, the reduction of the standard mass-spectrum relative to the studied one is 50 times. Natural Ni has five isotopes (see Table 8.8). Nuclides ^{58}Ni and ^{60}Ni possess high contents in the nature, the

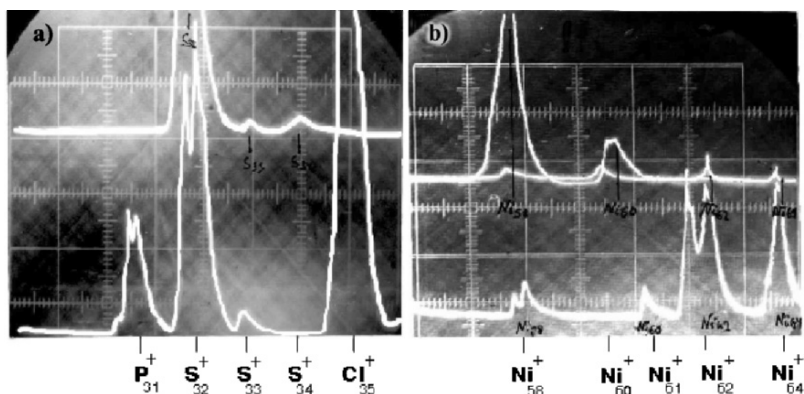


Fig. 8.3. Deviation of the isotope composition S (a) and Ni (b) contained in explosion products from their natural isotope composition.

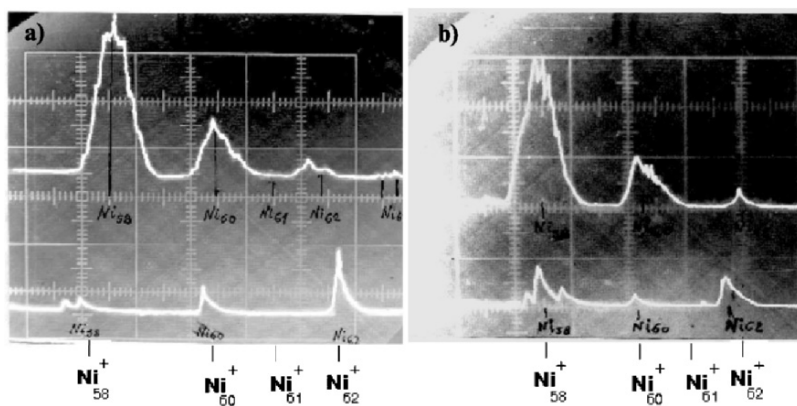


Fig. 8.4. Anomalies of the isotope composition of Ni contained in the explosion products of gold (a) and copper (b) targets.

first being the basic one. Likewise, the contents of nuclides ^{61}Ni and ^{64}Ni are about 1%, and the content of nuclide ^{62}Ni exceeds them approximately by 3.5 times. Comparing the upper (standard Ni) and lower (a specimen) parts of Fig. 8.3, b, we note the extremely high content of nuclides ^{61}Ni , ^{62}Ni , and ^{64}Ni in the explosion products of a platinum target and the absence of nuclide ^{60}Ni .

In the mass-spectrum registered on the explosion products of a gold target and given in Fig. 8.4, a, the mass-peaks of nuclides ^{60}Ni and ^{62}Ni have amplitudes, respectively, by two and five times more than that of the mass-peak of the basic nuclide ^{58}Ni . We note that their amplitudes should be, respectively, 2.6 and 19 times less than the amplitude of its peak in the case of the natural abundance.

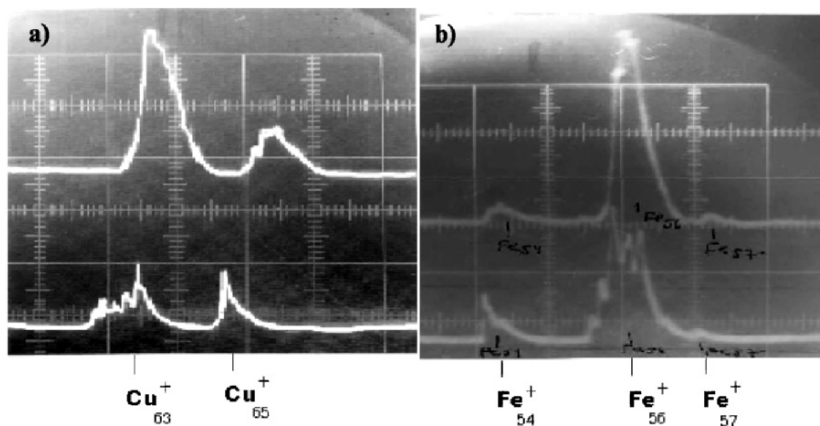


Fig. 8.5. Anomalies of the isotope composition of Cu (a) and Fe (b) contained in the explosion products of platinum and copper targets, respectively.

The mass-spectrum registered on the explosion products of a copper target and given in Fig. 8.4, b, illustrates one more case of anomalies of the isotope composition of Ni. Here, nuclides ^{58}Ni and ^{62}Ni have the approximately equal amplitudes of mass-peaks. But, in natural Ni, the amplitude of the mass-peak of the latter must be by 19 times less than that of the former.

In the upper and lower parts of Fig. 8.5, a, we give, respectively, the reference fragment of the mass-spectrum of a Cu standard and the corresponding fragment of the mass-spectrum of Cu registered on the explosion products of a platinum target on Cu accumulating screen No. 3013. The reduction of the reference mass-spectrum relative to the studied one is 20 times. Natural Cu has two nuclides: ^{63}Cu and ^{65}Cu . The first is basic, and its content in natural Cu exceeds the content of the second nuclide by more than two times (see Table 8.8). However, the mass-spectrum registered on the investigated specimen and given in Fig. 8.5, a, shows that these nuclides have the approximately equal amplitudes of mass-peaks. That is, the explosion products of a platinum target have the content of nuclide ^{65}Cu which is excessive by two times.

Fig. 8.5, b illustrates the anomaly of the isotope composition of Fe. The fragment of the mass-spectrum presented in its lower part was registered on the explosion products of a Cu target on Cu accumulating screen No. 2569. In this case, the reduction of the fragment of the reference mass-spectrum relative to it is 20 times. In the nature, the chemical element Fe is represented by four nuclides (see Table 8.8). The basic nuclide of Fe is ^{56}Fe . Its content in the nature exceeds those of nuclides ^{54}Fe , ^{57}Fe , and ^{58}Fe , respectively, by 16, 42, and 330 times. In the mass-spectrum registered

on the investigated specimen, we see the clearly excessive content of nuclide ^{54}Fe . Here, the amplitude of its mass-peak is only by 2.5 times less than that of the basic nuclide of Fe, rather than by 16 times as in the case of the natural distribution of nuclides.

Thus, we may conclude that the anomalies of the isotope composition for such chemical elements as S, Cl, Ca, Fe, Ni, Cu, and others registered by laser mass-spectrometry, as well as the establishment of the fact of the absence of a separation of isotopes upon the dispersion of the target substance, indicate that the explosion of a target substance is accompanied by its nuclear regeneration.

Isotope Composition of Chemical Elements Contained in Explosion Products and Being Minor Components of the Initial Materials of Targets and Accumulating screens. The determination of the isotope composition of minor chemical elements in explosion products was carried out by glow-discharge mass-spectrometry (VG 9000, VG Elemental, UK) and secondary-ion mass-spectrometry (IMS 4f, CAMECA, France). In the first case, we used the same mode of operation of the setup and the procedure of registration of spectra as upon studying the isotope composition of the basic chemical elements of target materials in explosion products. We also applied one of two previously used schemes of analysis, in which the 5-mm spot of a region analyzed with a mass-spectrometer was positioned directly on the surface of a layer of target explosion products deposited on the working surface of an accumulating screen (see Fig. 8.1, b). Even for minor chemical elements of the materials of targets, this scheme of analysis ensured a high signal upon the registration of the ion currents of their nuclides. However, due to a limited thickness of the layer of explosion products, it did not allow us to carry out the analysis of the isotope composition of chemical elements contained in them in a sufficiently wide range of masses. In fact, we were able to correctly perform the isotope analysis at an etching rate of about 0.1 to 0.2 $\mu\text{m}/\text{min}$ and a thickness of the layer of explosion products of 3 to 4 μm only for several chemical elements.

Therefore, upon the analysis of the isotope composition according to the mentioned scheme, we studied an incomplete range of the masses of nuclides belonging to the composition of target explosion products. We gave preference usually to those chemical elements which were contained in explosion products in significantly greater amounts, than in the initial materials of a target and an accumulating screen. We also paid attention to the vicinity of the masses of nuclides of iron (24 to 71 a.m.u.). This mass range was of interest for the study by virtue of that the nuclides belonging to it possess a maximum binding energy per nucleon and, hence, are the most

probable candidates to representatives of the nucleosynthesis products. To the priority group, we referred also chemical elements, whose anomalies of the isotope composition were already registered earlier by other methods.

The determination of the isotope composition of the analyzed chemical elements was performed only on those sections of a layer of target explosion products which were homogeneous by composition and were separated by means of the control over and comparison of the ion currents of reference nuclides in the subsequent cycles of scanning over the collection of windows chosen for the registration upon the continuous etching of the studied layer. The mentioned procedure of separation of homogeneous sections from a layer of explosion products is minutely described in the section devoted to the analysis of the isotope composition of basic chemical elements of the initial material of targets in explosion products. We also note that, upon the use of the discussed scheme of analysis, the ion currents of the studied nuclides fell always in the working range of a photoelectronic multiplier upon the registration of currents due to their low concentrations in explosion products. In other words, the measurements did not involve any errors of the determination of the isotope composition related to the insufficient accuracy of the consistency of detectors used in the registration of ion currents.

The second scheme of analysis upon the determination of the isotope composition of target explosion products by glow-discharge mass-spectrometry was basically different from the first one only by the type of specimens under investigation. In this case, we studied specially manufactured specimens of the "sandwich" type, being an assembly designed in the form of a stack made of several tens of screens closely adjacent one to another with explosion products (see Fig. 8.24), rather than individual accumulating screens with deposited explosion products. Such a stack was cut in halves, and the formed surfaces consisting of the cut ends of accumulating screens were analyzed. Just on such a surface was located a 5-mm spot of the region analyzed by a mass-spectrometer.

The significant merit of specimens of the above-mentioned construction is their homogeneity in the direction of analysis. Their construction ensured a good averaging of inhomogeneities of both the composition of individual accumulating screens and the dispersion of the results of various experiments. This circumstance allowed us to remove any limitations on both the duration of a measuring experiment and the size of a collection of windows upon the registration of the ion currents of studied nuclides, i.e., we were able to register all the chemical elements contained in explosion products without limitations.

However, we should like to indicate the fact that the usage of specimens of the "sandwich" type for the determination of the isotope

composition of target explosion products has one troubling drawback in addition to the mentioned advantages. The large ratio of the cross-section area of the substrate (accumulating screens) to the cross-section area of layers of explosion products on a “sandwich” leads to that the anomalies of the isotope composition defined by the contribution of the studied layers have to be measured against a significant background of the natural distribution of isotopes, which is conditioned by the contribution from the substrate. A decrease in the level of the effective signal is the price which should be paid for the homogeneity of the specimens under study. It is clear that the anomalies of the isotope composition of chemical elements which are registered on “sandwiches” are considerably less than the primary anomalies of the isotope composition of chemical elements contained in explosion products.

Finally, we note that the mass composition of the glow-discharge plasma was calculated always for both the schemes of analysis for the clarification of, and to account for, a possible superposition of the mass-peaks of molecular ions on the measured mass-peaks. We also considered the interferences of the mass-peaks of the studied nuclides with all mass-peaks of clusters possessing a significant intensity. Such clusters usually contain various combinations of the nuclides of basic chemical elements of a target material and an accumulating screen, argon, and associated gases. In the presence of such interferences, the analyzed mass-peaks were not considered.

We now illustrate some anomalies of the isotope composition of minor chemical elements of a target material registered in explosion products by glow-discharge mass-spectrometry. In Table 8.9, we present the results of the determination of the isotope composition of a number of minor chemical elements in target explosion products. The measurements were carried out on a specimen of the “sandwich” type collected from copper accumulating screens. We will comment on the form of the presentation of data in Table 8.9. In the cases where the mass-peaks of some nuclides of the analyzed chemical elements underwent interferences, their symbols were marked by two asterisks, and we put a dash instead of a datum to the corresponding cell. In these cases, the symbols of the remaining nuclides of a chemical element were marked by one asterisk, and the content of each nuclide was normalized on the sum of contents of all other analyzed nuclides of the chemical element under consideration.

It is seen from the given data that the deviation of the content of nuclides of some chemical elements in target explosion products from their natural abundance can reach several tens (^{13}C , ^{33}S , ^{34}S , ^{46}Ti , ^{49}Ti , ^{50}Ti , ^{67}Zn , and ^{68}Zn) and even several hundreds of percents (^{58}Fe). We note also that the large anomalies of the isotope composition are usually typical of

Table 8.9. Results of the determination of the isotope composition of some minor chemical elements in target explosion products (target material – Pb) on a specimen of the “sandwich” type (copper accumulating screens).

Nuclide	Conc., %	δ , %	Nat. [†] , %
¹² C	99.13	0.23	98.9
¹³ C	0.87	–20.9	1.1
⁶³ Cu	70.69	2.2	69.17
⁶⁵ Cu	29.30	–4.9	30.83
³² S*	93.83	–1.27	95.04
³³ S*	0.66	–12.02	0.75
³⁴ S*	5.51	30.85	4.21
³⁶ S**	–	–	–
⁴⁶ Ti	9.39	17.4	8.0
⁴⁷ Ti	7.88	7.9	7.3
⁴⁸ Ti	72.93	–1.18	73.8
⁴⁹ Ti	6.24	13.5	5.5
⁵⁰ Ti	3.55	–34.0	5.4
⁵⁴ Fe**	–	–	–
⁵⁶ Fe*	96.53	–0.86	97.37
⁵⁷ Fe*	2.32	–0.81	2.34
⁵⁸ Fe*	1.16	289.69	0.30
⁶⁴ Zn	52.15	7.3	48.6
⁶⁶ Zn	27.66	–0.9	27.9
⁶⁷ Zn	3.53	–13.9	4.1
⁶⁸ Zn	16.05	–14.6	18.8
⁷⁰ Zn	0.61	1.7	0.6

* concentrations are normed on the sum of the concentrations of analyzed nuclides.

** there occurs the interference of mass-peaks.

† natural abundance (see Refs. 100, 101).

nuclides with low abundance in nature. However, not less significant and convincing are the anomalies of the isotope composition equal to several percents for the basic nuclides of chemical elements (⁶⁴Zn) or for nuclides with high abundance (⁶⁵Cu).

We also would point out that the anomalies of the isotope composition of chemical elements registered on “sandwiches” are reduced to a considerable extent as compared to their anomalies in target explosion products due to their strong dissolution into the substrate material (the material of

accumulating screens) which have the natural abundance of nuclides. In view of this point and the fact that the basic chemical element of accumulating screens in the discussed experiment was copper, the deviation of the content of nuclide ^{65}Cu by 4.9% from its natural abundance looks rather surprising. On the one hand, the indicated anomaly exceeds the measurement error by almost one order of magnitude. On the other hand, it is obvious that the amount of the substance of target explosion products transferred on the accumulating screen is too small in order to change the isotope composition of a huge mass of copper contained in the accumulating screen to such a degree as a result of the mixing. A probable explanation of this fact can be apparently related only to an undefined but evident nuclear regeneration of the screen material as a result of the explosion of a target.

We now consider the anomalies of the isotope composition registered on individual accumulating screens. In Table 8.10, we give the results of the determination of the isotope composition of sulfur in target explosion products deposited on copper accumulating screens. The mass-peaks of all nuclides of sulfur, except for ^{36}S , are reliably registered. The concentrations of nuclides ^{32}S and ^{34}S in explosion products only slightly differ from their natural abundance. The deviations of the concentration of nuclide ^{33}S from its natural content are much greater and reach 30% to 40% by shifting to the lower level in all the cases. As for nuclide ^{36}S , its content in explosion products and in nature is extremely low, and its mass-peak underwent a strong interference with the mass-peak of ^{36}Ar .

The results of the analysis of the isotope composition of titanium are given in Table 8.11. The mass-peaks of all nuclides of Ti are reliably registered. By estimating the mass-peak amplitude of nuclide ^{50}Ti , we took into account the influence of a insignificant contribution from nuclide ^{50}Cr . It is seen from the presented data that the deviations of the content of isotopes of titanium in explosion products from their natural abundance for ^{46}Ti , ^{47}Ti , ^{48}Ti , ^{49}Ti , and ^{50}Ti are significant and reach the level of 68.5%, 31.1%, -15.37%, 23.45%, and 109.26%, respectively. We note that the indicated deviations vary from experiment to experiment in a wide range. However, they bear a, more or less, constant character in that the basic nuclide ^{48}Ti has lower content in explosion products, than its natural abundance, and the remaining nuclides of titanium have, vice versa, higher content.

The results of the analysis of the isotope compositions of Cr and Fe in explosion products are presented in Table 8.12. Commenting on these results, we note, first of all, that the weak mass-peak of nuclide ^{54}Cr is superimposed on the mass-peak of nuclide ^{54}Fe with a comparable intensity, and the weak mass-peak of nuclide ^{58}Fe is superimposed on the intense mass-peak of nuclide ^{58}Ni . These circumstances lead to the absence of the

Table 8.10. Results of the determination of the isotope composition of sulfur in lead target explosion products on copper accumulating screens.

Sample	Nuclide	Conc., %	δ , %	Nat. [†] , %
84 (edge)	³² S*	95.15	0.12	95.04
	³³ S*	0.52	−30.68	0.75
	³⁴ S*	4.33	2.83	4.21
	³⁶ S**	—	—	—
84 (center)	³² S*	95.39	0.37	95.04
	³³ S*	0.53	−29.35	0.75
	³⁴ S*	4.08	−3.11	4.21
	³⁶ S**	—	—	—
84 (center)	³² S*	94.98	−0.06	95.04
	³³ S*	0.46	−38.68	0.75
	³⁴ S*	4.56	8.29	4.21
	³⁶ S**	—	—	—
88 (center)	³² S*	94.98	−0.06	95.04
	³³ S*	0.62	−17.62	0.75
	³⁴ S*	4.40	4.48	4.21
	³⁶ S**	—	—	—
90 (center)	³² S*	94.95	−0.09	95.04
	³³ S*	0.45	−40.01	0.75
	³⁴ S*	4.60	9.24	4.21
	³⁶ S**	—	—	—

* concentrations are normed on the sum of the concentrations of analyzed nuclides.

** there occurs the interference of mass-peaks.

† natural abundance (see Refs. 100, 101).

possibility of properly registering the ion currents of isotopes ⁵⁴Cr, ⁵⁴Fe, and ⁵⁸Fe. By virtue of the above-written, we give no data on these nuclides in Table 8.12, and the data on the remaining nuclides of Cr and Fe are presented to be normalized on the sum of the contents of the analyzed nuclides. It follows from the data of Table 8.12 that isotopes ⁵⁰Cr and ⁵⁷Fe in explosion products reveal a somewhat increased content as compared to their natural abundance, whereas isotope ⁵³Cr shows a somewhat decreased content. Moreover, the indicated tendency has a quite stable character.

The results of the analysis of the isotope composition of nickel and zinc in lead target explosion products are given in Table 8.13. For nickel, the

Table 8.11. Results of the determination of the isotope composition of titanium in lead target explosion products on copper accumulating screens.

Sample	Nuclide	Conc., %	δ , %	Nat. [†] , %
84 (edge)	⁴⁶ Ti	10.55	31.88	8.00
	⁴⁷ Ti	7.78	6.58	7.30
	⁴⁸ Ti	69.73	-5.51	73.80
	⁴⁹ Ti	4.91	-10.73	5.50
	⁵⁰ Ti	7.03	30.19	5.40
84 (center)	⁴⁶ Ti	9.03	12.88	8.00
	⁴⁷ Ti	7.30	0.00	7.30
	⁴⁸ Ti	72.29	-2.05	73.80
	⁴⁹ Ti	6.79	23.45	5.50
	⁵⁰ Ti	4.60	-14.81	5.40
84 (center)	⁴⁶ Ti	13.48	68.50	8.00
	⁴⁷ Ti	9.57	31.10	7.30
	⁴⁸ Ti	62.46	-15.37	73.80
	⁴⁹ Ti	6.24	13.45	5.50
	⁵⁰ Ti	8.23	52.41	5.40
88 (center)	⁴⁶ Ti	7.28	-9.00	8.00
	⁴⁷ Ti	6.80	-6.85	7.30
	⁴⁸ Ti	73.26	-0.73	73.80
	⁴⁹ Ti	5.85	6.36	5.50
	⁵⁰ Ti	6.80	25.93	5.40
90 (center)	⁴⁶ Ti	12.18	52.25	8.00
	⁴⁷ Ti	7.92	8.49	7.30
	⁴⁸ Ti	62.72	-15.01	73.80
	⁴⁹ Ti	5.86	6.55	5.50
	⁵⁰ Ti	11.3	109.26	5.40

[†] natural abundance (see Refs. 100, 101).

ion currents of nuclides ⁶⁰Ni, ⁶¹Ni, and ⁶²Ni are measured reliably, the intense mass-peak of nuclide ⁵⁸Ni interferes with the low-intensity mass-peak of nuclide ⁵⁸Fe, and the low-intensity mass-peak of nuclide ⁶⁴Ni interferes with the intense peak of the basic nuclide ⁶⁴Zn. By virtue of the above-written, the data concerning the last isotope of nickel in Table 8.13 are absent. As for zinc, only the mass-peaks of its basic isotopes ⁶⁴Zn and ⁶⁶Zn touch the “tails” (the Cu substrate) of lines ⁶³Cu¹H¹ and ⁶⁵Cu¹H¹, respectively, whose contribution can be easily taken into account. Analyzing the given data, we

Table 8.12. Results of the determination of the isotope composition of Cr and Fe in lead target explosion products on copper accumulating screens.

Sample	Nuclide	Conc., %	δ , %	Nat. [†] , %
84 (edge)	⁵⁰ Cr*	3.80	−14.59	4.45
	⁵² Cr*	87.37	1.81	85.82
	⁵³ Cr*	8.82	−9.31	9.73
	⁵⁴ Cr**	—	—	—
84 (center)	⁵⁰ Cr*	5.54	24.51	4.45
	⁵² Cr*	86.58	0.89	85.82
	⁵³ Cr*	7.88	−19.04	9.73
	⁵⁴ Cr**	—	—	—
84 (center)	⁵⁰ Cr*	6.55	47.20	4.45
	⁵² Cr*	82.92	−3.38	85.82
	⁵³ Cr*	10.53	8.22	9.73
	⁵⁴ Cr**	—	—	—
88 (center)	⁵⁰ Cr*	4.51	1.24	4.45
	⁵² Cr*	91.48	6.59	85.82
	⁵³ Cr*	4.02	−58.72	9.73
	⁵⁴ Cr**	—	—	—
90 (center)	⁵⁰ Cr*	—	—	4.45
	⁵² Cr*	—	—	85.82
	⁵³ Cr*	—	—	9.73
	⁵⁴ Cr**	—	—	—
84 (edge)	⁵⁴ Fe**	—	—	—
	⁵⁶ Fe*	97.41	−0.25	97.66
	⁵⁷ Fe*	2.59	10.62	2.34
	⁵⁸ Fe**	—	—	—
84 (center)	⁵⁴ Fe**	—	—	—
	⁵⁶ Fe*	97.61	−0.05	97.66
	⁵⁷ Fe*	2.39	2.09	2.34
	⁵⁸ Fe**	—	—	—
84 (center)	⁵⁴ Fe**	—	—	—
	⁵⁶ Fe*	97.26	−0.41	97.66
	⁵⁷ Fe*	2.74	17.09	2.34
	⁵⁸ Fe**	—	—	—

	$^{54}\text{Fe}^{**}$	—	—	—
88	$^{56}\text{Fe}^*$	97.54	−0.12	97.66
(center)	$^{57}\text{Fe}^*$	2.46	4.94	2.34
	$^{58}\text{Fe}^{**}$	—	—	—
	$^{54}\text{Fe}^{**}$	—	—	—
90	$^{56}\text{Fe}^*$	97.29	−0.38	97.66
(center)	$^{57}\text{Fe}^*$	2.71	15.87	2.34
	$^{58}\text{Fe}^{**}$	—	—	—

* concentrations are normed on the sum of the concentrations of analyzed nuclides.

** there occurs the interference of mass-peaks.

† natural abundance (see Refs. 100, 101).

indicate the contents of isotopes ^{61}Ni and ^{62}Ni in target explosion products which considerably differ from the natural abundance. As for the isotope composition of zinc, it is considerably shifted to the side of nuclides with lesser content of neutrons.

Finally, we present the results of the analysis of the isotope composition of lead in the explosion products of a silver target on a Cu accumulating screen (see Table 8.14). All its isotopes are in a good analytic situation and undergo no considerable interferences with other masses. It is seen from the presented data that Pb contained in the explosion products of a target is significantly different from natural lead by the isotope composition. First of all, we may note a significant increase in the content of nuclides ^{204}Pb and ^{206}Pb and a decrease in the content of the basic nuclide Pb^{208} in explosion products. As for nuclide ^{207}Pb , the deviations of its content are insignificant. In other words, we observe a remarkable shift of its isotope composition to the side of nuclides with lesser content of neutrons like in the case where Pb was the basic chemical element of the target material (see Table 8.6).

As noted above, the determination of the isotope composition of minor chemical elements in the explosion products was also performed by secondary-ion mass-spectroscopy (IMS 4f, CAMECA, France). As ions of a primary beam, we used mostly ions O_2^+ and Cs^+ and sometimes ions N_2^+ and Xe^+ , their energies being varied from 10 to 15 keV. Depending on the tasks and purposes of the analysis, the primary beam current was from 1 to 100 nA. But, in most cases, it was in the range 4 to 6 nA. The size of an ion probe raster was taken to be $500 \times 500 \mu\text{m}$. These conditions provided low rates of the spraying of the surface of a studied specimen and, hence, the mode of analysis close to the static one. This allowed us to carry out

Table 8.13. Results of the determination of the isotope composition of nickel and zinc in lead target explosion products on copper accumulating screens.

Sample	Nuclide	Conc., %	δ , %	Nat. [†] , %
84 (edge)	⁵⁸ Ni*	68.84	−0.08	68.90
	⁶⁰ Ni*	26.36	0.08	26.34
	⁶¹ Ni*	1.23	7.86	1.14
	⁶² Ni*	3.57	−1.46	3.62
	⁶⁴ Ni**	—	—	—
84 (center)	⁵⁸ Ni*	71.80	4.21	68.90
	⁶⁰ Ni*	24.24	−7.96	26.34
	⁶¹ Ni*	1.27	11.38	1.14
	⁶² Ni*	2.69	−25.74	3.62
	⁶⁴ Ni**	—	—	—
84 (center)	⁵⁸ Ni*	68.93	0.05	68.90
	⁶⁰ Ni*	24.33	−7.63	26.34
	⁶¹ Ni*	2.40	110.46	1.14
	⁶² Ni*	4.34	19.79	3.62
	⁶⁴ Ni**	—	—	—
88 (center)	⁵⁸ Ni*	67.99	−1.32	68.90
	⁶⁰ Ni*	27.23	3.38	26.34
	⁶¹ Ni*	1.25	9.61	1.14
	⁶² Ni*	3.53	−2.57	3.62
	⁶⁴ Ni**	—	—	—
90 (center)	⁵⁸ Ni*	68.10	−1.16	68.90
	⁶⁰ Ni*	26.07	−1.01	26.34
	⁶¹ Ni*	1.57	37.69	1.14
	⁶² Ni*	4.26	17.59	3.62
	⁶⁴ Ni**	—	—	—
84 (edge)	⁶⁴ Zn	50.89	4.70	48.60
	⁶⁶ Zn	28.22	1.14	27.90
	⁶⁷ Zn	4.06	−0.90	4.10
	⁶⁸ Zn	16.27	−13.47	18.80
	⁷⁰ Zn	0.57	−5.81	0.60
84 (center)	⁶⁴ Zn	50.74	4.40	48.60
	⁶⁶ Zn	27.94	0.15	27.90
	⁶⁷ Zn	3.58	−12.71	4.10

	^{68}Zn	17.16	-8.73	18.80
	^{70}Zn	0.58	-3.18	0.60
84 (center)	^{64}Zn	51.17	5.30	48.60
	^{66}Zn	28.21	1.12	27.90
	^{67}Zn	3.19	-22.15	4.10
	^{68}Zn	16.83	-10.49	18.80
	^{70}Zn	0.59	-1.04	0.60
88 (center)	^{64}Zn	50.21	3.31	48.60
	^{66}Zn	28.01	0.39	27.90
	^{67}Zn	4.00	-2.44	4.10
	^{68}Zn	17.18	-8.62	18.80
	^{70}Zn	0.60	0.00	0.60
90 (center)	^{64}Zn	50.86	4.65	48.60
	^{66}Zn	28.06	0.58	27.90
	^{67}Zn	3.91	-4.70	4.10
	^{68}Zn	16.57	-11.86	18.80
	^{70}Zn	0.60	0.00	0.60

* concentrations are normed on the sum of the concentrations of analyzed nuclides.

** there occurs the interference of mass-peaks.

† natural abundance (see Refs. 100, 101).

Table 8.14. Results of the determination of the isotope composition of Pb in explosion products on specimen No. 109 in the experiment with a silver target and a copper accumulating screen.

Sample	Nuclide	Conc., %	δ , %	Nat. [†] , %
	^{204}Pb	1.59	13.57	1.4
109	^{206}Pb	26.60	10.37	24.1
(center)	^{207}Pb	23.12	4.62	22.1
	^{208}Pb	48.69	-7.08	52.4

† natural abundance (see Refs. 100, 101).

long-term measurements without significant change in the composition of the studied area of the specimen surface. That is, we did not take care of the selection of the homogeneous layers of target explosion products for the analysis. Large sizes of the ion probe raster were chosen with the purpose

both to increase the collection area of secondary ions and to compensate the inevitable loss in the sensitivity of the method at low currents of the primary ion beam.

The used ion optics allowed us to construct the image of a studied area of the surface in the form of a 2D distribution of analyzed mass with spatial resolution to be better than $1\text{ }\mu\text{m}$. In this case, to get high-grade images, we used long-term exposures (in some cases, up to 20 min and more). The analyzed range of a magnetic separator of masses was 1 to 480 a.m.u. In the used working modes, the secondary-ion optics ensured the spectral resolution $M/\Delta M$ not lower than 6000 upon the conservation of the sensitivity to be sufficient for the solution of the posed research problems. To register the ion currents, the setup included a photoelectronic multiplier (the range of small currents) and a Faraday cup (the range of large currents). The analysis of the isotope composition of target explosion products was performed in the range of registration of the currents of a photoelectronic multiplier. That is, the measurements had no errors of the determination of the isotope composition related to an insufficient accuracy of the consistency of the detectors used in the registration of ion currents.

By secondary-ion mass-spectrometry, we analyzed the isotope composition of a layer of solid products of the explosion of a target. This layer was deposited on an accumulating screen of the standard form and size (see Fig. 8.1, b). As the initial material of an accumulating screen, we usually took chemically pure Cu and sometimes Ag, Ta, Au, Pt, etc. In the production of targets, we used a wider collection of materials, such as, light, medium, and heavy pure metals with atomic masses in the interval from 9 to 209.

Upon the study of the isotope composition of explosion products, the microprobe raster (the analyzed area) was positioned directly on the surface of the layer of target explosion products deposited on an accumulating screen. In this case, the direction of the analysis was normal to the working surface of the accumulating screen. With the purpose of providing the representativeness of results of the determination of the isotope composition, the choice of the positions of analyzed areas on the specimens under study was realized according to axial symmetry criteria. First, we chose several most typical radial directions on accumulating screens and then carried out the analysis at 7 to 10 points positioned more or less uniformly along each direction.

Finally, we consider the very procedure of analysis of the mass-spectra registered by secondary-ion mass-spectrometry, because its correctness and the adequacy of the decoding of mass-spectra define directly the reliability of the derived results (in our case, the isotope composition of explosion products). In the general case, the mass-spectrum of secondary ions is quite

complicated, which is reflected in the presence of a huge number of peaks in mass-spectra and in the frequent cases of their superpositions one onto another (interferences). In particular, the spectrum contains the mass-peaks of monoatomic (single-charge and multiply charged) ions of the basic components of a material and its admixtures; ions of two-, three-, and many-atom clusters. The last ones are formed on the base of chemical elements with significant concentrations, the ions of active elements of the primary beam and the ions of chemical compounds, including the compounds of the basic components of the material with oxygen, nitrogen, hydrogen, etc.

The main difficulties of the correct decoding of such mass-spectra, i.e., the difficulties of the identification of the peaks corresponding to one-atom ions, ions of clusters, and ions of chemical compounds, are defined by the occurrence of their interference, rather than by their huge amount. The interferences are mainly created by the mass-peaks of complex and multiply charged ions. In this case, the latter present considerably less difficulties by virtue of the fact that the amplitudes of their peaks are extremely low. Usually the number of two-charge ions is observed to exist by two to three orders of magnitude less than that of one-charge ions, and so on (see Refs. 99, 102–104). To successfully decode complex mass-spectra, it is necessary, first of all, to register them with sufficiently high resolution (at least, with $M/\Delta M$ not less than 1000) in order to separate the most part of doublets and triplets contained in them and, by this, to considerably decrease the number of interferences. The last circumstance is quite important by virtue of the fact that a complex mass-spectrum is not, strictly saying, a linear superposition of the mass-spectra of appropriate pure chemical elements. One of the reasons of this fact is the great difference in the yield coefficients of secondary-ion emission of chemical elements (see Ref. 102). In this connection, we recall that the spectral resolution of the mass-spectrometer used by us is reliably satisfied the condition formulated above.

Upon the analysis of the mass-spectra registered on target explosion products, the decoding began from the identification of peaks, whose mass numbers correspond to the most abundant isotopes of chemical elements contained in them. The adequacy of the decoding for many-isotope elements was controlled by comparison of the ratios of the amplitudes of relevant peaks with those characteristic of the natural abundance of the isotopes of these elements. On the second stage of the analysis, we identified the multiply charged (as usual, two-charge) ions of the main components. In this case, we took into account the circumstance that their mass numbers in a mass-spectrum, being actually the ratios M/q , are twice less than the mass numbers of the corresponding one-charge ions. We began the identification of the ions of clusters from a search for the two-atom combinations of

atoms of the main components and carried it out with regard to the equality of the probabilities of the combinations of various isotopes of a given element. We note that, as a rule, the number of complex ions sharply decreases with increase in the number of atoms belonging to them. In the next stage, we identified the peaks of the mass-spectrum which correspond to the secondary ions of oxides M_xO_y , hydrides M_xH_y , hydroxides $M_x(OH)_y$, and hybrid molecules M_xN_y , which are typical of mass-spectrometry.

After the execution of a decoding of the mass-spectrum by the procedure described above, we separated the groups of mass-peaks in it which are related to the isotopes of a single chemical element. Then, for each group, we checked the correspondence of isotope ratios to their natural abundance. If we found a disagreement, the mass-peaks with anomalous values of the amplitudes were tested for the presence of their interferences with the mass-peaks of complex and multiply charged ions. The mentioned testing procedure of the mass-peaks with anomalous intensities consisted of three stages. A deviation of the isotope composition of the chemical element, to which the isotope corresponding to the tested mass-peak belongs, was considered to be registered only if all three stages of the testing procedure were successfully passed. In the opposite case where the performed testing showed interferences, the studied mass-peak and the corresponding isotope were excluded from the consideration.

Below, we describe the essence of the testing procedure for mass-peaks for the presence of their interferences with the mass-peaks of complex and multiply charged ions. Its first stage consisted in the repeated registration of a fragment of the mass-spectrum containing the studied peak in the offset mode. In most cases, the mentioned mode efficiently suppressed the mass-peaks of cluster ions. Its principle is based on the fact that secondary atomic ions possess usually a wider distribution over a specific energy range than secondary complex ions (see Fig. 8.6).

In the normal working mode, the transmitting slit of the energy filter of a mass-spectrometer is in position I which defines the ratio of the amplitudes of registered mass-peaks (in the case under consideration, of the peaks of ions Si^+ , SiH^+ , and Si_2O^+). Switching on the offset mode shifts the transmitting slit of the energy filter of a mass-spectrometer by several tens of eV to the high-energy part of the energy distribution of secondary ions (see position II in Fig. 8.6) and transfers, in fact, the work of a mass-spectrometer to the mode of collection of only fast ions. This possibility is technically realized by means of the supply of a small cutting-off potential (usually up to 90 V) to the studied specimen. It follows from Fig. 8.6 that, in the last case, the ratio of the amplitudes of registered mass-peaks is sharply changed in favor of the mass-peak corresponding to the atomic ion Si^+ . That is, the

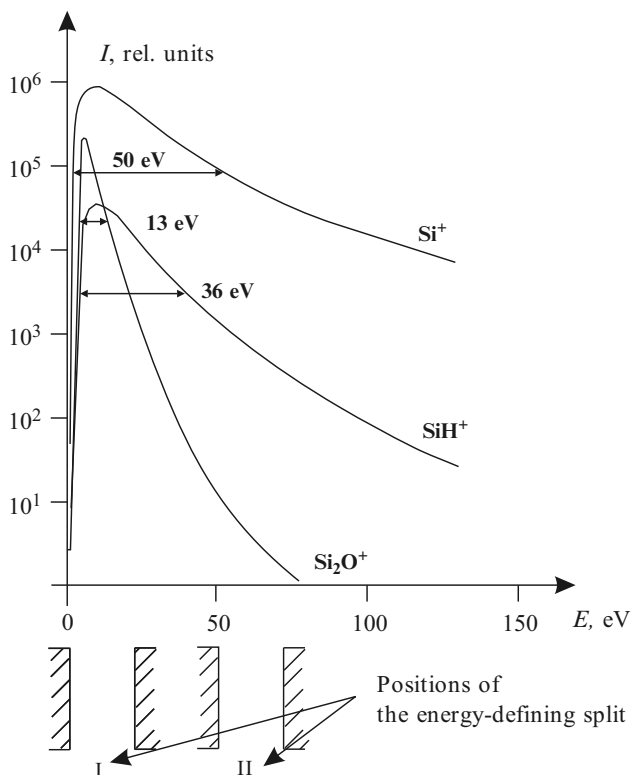


Fig. 8.6. Character of the energy distribution of secondary atomic (Si^+) and complex (SiH^+ and Si_2O^+) ions (see Ref. 104) and the scheme of a shift of the energy-defining split of a mass-spectrometer which illustrates the principle of realization of the offset mode.

supply of the offset voltage leads to the efficient discrimination of the mass-peaks of complex ions.

Thus, it follows from the above-presented that if the tested mass-peak preserves the anomalous value of its amplitude relative to the amplitudes of mass-peaks of other isotopes of the studied chemical element with increase in the offset voltage, then its interferences with the mass-peaks of clusters are, most likely, absent, and we are faced with a true deviation the isotope composition. In the opposite case, we may say with a great degree of probability that the anomalous value of its amplitude is caused by the interference.

The second stage of the procedure of testing the mass-peaks possessing the anomalous amplitude for the presence of interferences is extremely important in rather rare cases where the energy distributions of atomic and complex ions with identical nominal masses differ insignificantly, and the use of the offset mode does not solve the problem of recognition of the

interference of mass-peaks. In these cases, to identify the overlapping mass-peaks, one usually uses the recording mode with high resolution in mass for the analyzed fragment of a mass-spectrum. The gist of this trick consists in that the coincidence of nominal masses of atomic and complex ions does not mean the coincidence of their exact masses, and the recording of the spectrum fragment containing the tested peak in the mode with high resolution in mass allows usually one to register the difference of their exact mass numbers and, hence, to identify them, as well as to answer the question about the presence of the interference, if such has occurred.

On the third stage, the question about the belonging of the tested mass-peak to an atomic or complex ion is solved by means of the analysis of the image of the distribution of the corresponding mass on the studied area of the specimen surface. This method of identification of molecular complexes is based on the obvious fact that the image of the surface distribution of a molecular mass must coincide with the images of the surface distributions of masses of the nuclides composing it. If such coincidences are observed, then the analyzed mass-peak should be referred to a molecular ion. In the opposite case, it will correspond to an atomic ion. We note also that if the tested mass-peak belongs to an atomic ion, the image of the surface distribution of its mass must coincide with the images of the surface distributions of the masses of other isotopes of the same chemical element.

Finally, we indicate that our tremendous experience of using the described procedure of testing the mass-peaks for the presence of their interferences with the mass-peaks of complex and multiply charged ions (on the second and third stages) showed its high efficiency and reliability. It is worth noting that this procedure is also suitable in those cases where the atomic ion corresponding to the tested mass-peak is the single nuclide of an analyzed chemical element.

Below, we will demonstrate some anomalies of the isotope composition of minor chemical elements of the materials of targets which are registered in the products of explosions by secondary-ion mass-spectrometry. We start with the deviation of the isotope composition of Mg registered on a copper accumulating screen in the products of the explosion of a lead target. All mass-peaks of the nuclides of Mg were reliably identified and passed the procedure of testing for the absence of interferences with the mass-peaks of complex and multiply charged ions. The absence of interferences is confirmed by fragments of the spectra registered in the mode with high resolution. These fragments contain the mass-peaks of all nuclides of Mg which are presented in Fig. 8.7. In this case, their concentrations were 74.7% (^{24}Mg), 11.5% (^{25}Mg), and 13.8% (^{26}Mg), which corresponds to the deviations from their natural abundance by — 5.4% (^{24}Mg), 15.0% (^{25}Mg), and 25.3% (^{26}Mg).

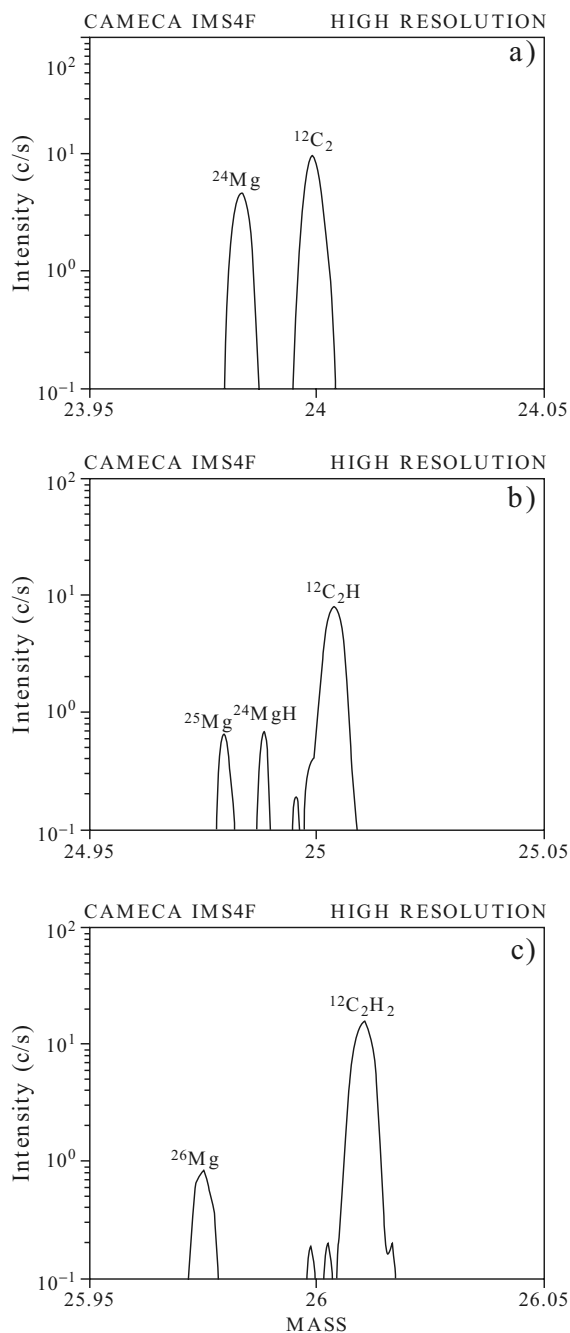


Fig. 8.7. Fragments of the mass-spectra of secondary ions which are registered in the mode with high resolution and contain the mass-peaks of nuclides ^{24}Mg (a), ^{25}Mg (b), and ^{26}Mg (c).

Other sufficiently bright and convincing examples of anomalies of the isotope composition were registered on specimen No. 96 (a Cu target and a Cu accumulating screen) upon the determination of the content of the nuclides of Li. Lithium is remarkable in that the mass-peaks of its nuclides have no significant interferences with the mass-peaks of complex and multiply charged ions. Indeed, clusters with mass numbers 6 and 7 must consist of nuclides lighter than Li, i.e., of H and He. However, the probability of the formation of such complexes is extremely low. As helium is an inert gas, and compounds formed with inert gasses, while not unknown, are typically highly unstable not to mention difficult to create. As for the interferences with multiply charged ions, their unique candidates are $^{12}\text{C}^{++}$ (for the mass number equal to 6) and $^{14}\text{N}^{++}$ (for the mass number equal to 7). However, if we take into account that, first, the yield coefficient of secondary-ion emission for carbon and nitrogen (upon the registration of positive secondary ions and with the primary beam of ions O_2^+) is four orders lower than that for lithium (see Ref. 99) and, second, the yield of their doubly charged ions is lower else by two orders of magnitude (see Ref. 99), it becomes obvious that their contributions to the intensity of the mass-peaks of ^6Li and ^7Li are negligible.

The above-presented is confirmed by fragments of the mass-spectra of secondary ions registered in the mode with high resolution and containing the mass-peaks of nuclides ^6Li and ^7Li (Fig. 8.8, a, b). As seen from Fig. 8.8, the mass number 6 or 7 corresponds to only one mass-peak which is spectrally clearly resolved. At the indicated point of analysis, the concentrations of the nuclides of Li were 6.6% (^6Li) and 93.4% (^7Li), which corresponds to the deviations from their natural abundance by -12.1% (^6Li) and 1.0% (^7Li).

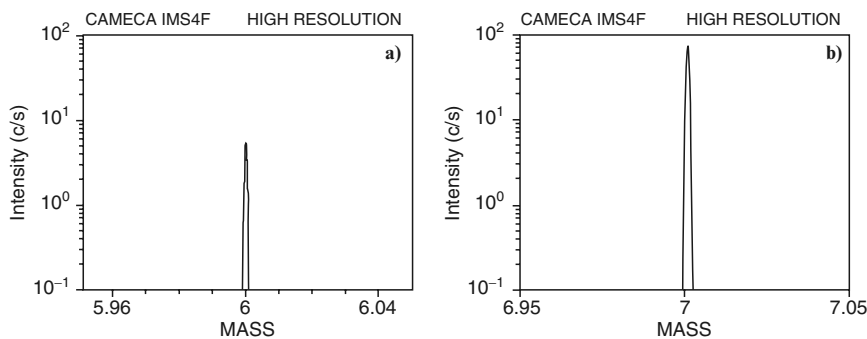


Fig. 8.8. Fragments of the mass-spectra of secondary ions which are registered in the mode with high resolution and contain the mass-peaks of nuclides ^6Li (a) and ^7Li (b), specimen No. 96, point 1.

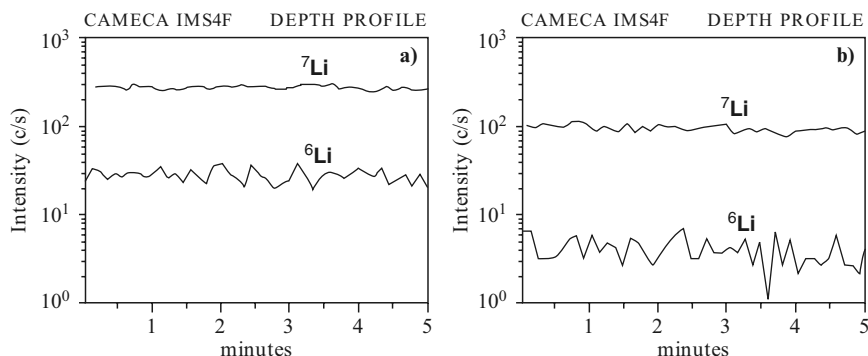


Fig. 8.9. Dependences of the intensities of the ion currents of nuclides ^6Li and ^7Li on time at point 2 (a) and point 3 (b), specimen No. 96.

On the same specimen, we also discovered two small regions with the aggregates of Li, in which we determined its isotope composition and found various anomalies. The results of these measurements are presented in Fig. 8.9 in the form of the temporal dependences of the intensities of the ion currents of both nuclides of Li. The recalculation of the ion currents of nuclides into their concentrations gave the following results: 4.2% ^6Li and 95.8% ^7Li at point 2 and 6.4% ^6Li and 93.6% ^7Li at point 3. The anomalies registered here bear the same character like that for point 1, and their values are -44.0% for ^6Li and 3.6% for ^7Li at point 2 and -14.7% for ^6Li and 1.2% for ^7Li at point 3, respectively.

One of the examples of anomalies of the isotope composition of palladium and barium is given in Fig. 8.10. It is represented in the clear form by means of the comparison of a fragment of the mass-spectrum of Pd and Ba registered on products of the explosion of a Cu target with the corresponding fragment of the standard mass-spectrum with the natural distribution of their nuclides. The latter fragment is positioned to the right from the former. The presented data indicate clearly that the isotope composition of Ba is shifted to the side of nuclides with lesser mass. As for the anomaly of the isotope composition of Pd, its character is more complex.

The next example presents the anomalies of the isotope composition found upon the study of the explosion products of a tantalum target which were deposited on the Cu accumulating screen No. 10623. The explosion products of a target on accumulating screens look frequently as drops elongated and introduced into a screen and as splashes dispersed from the screen center to its periphery. Sometimes, the dispersed splashes of a target material recoil from an accumulating screen by forming dents and craters. In Fig. 8.11, a, we show the image of one of such dents which is obtained in secondary ions $^{23}\text{Na}^+$. The dent was located on the accumulating screen

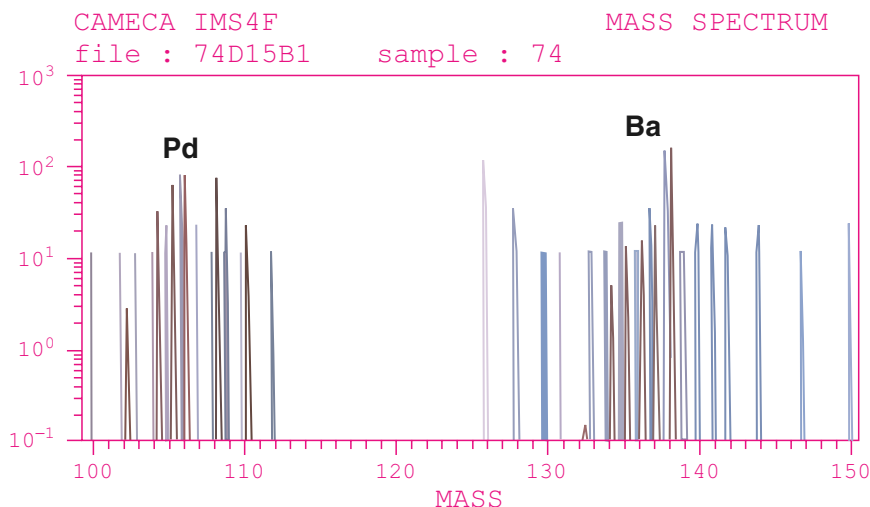


Fig. 8.10. Comparison of a fragment of the mass-spectrum of Pd and Ba registered on the target explosion products on specimen No. 74 (the filled peaks positioned to the left from the standard spectrum) with the corresponding fragment of a mass-spectrum with the natural distribution of their nuclides (the unfilled peaks).

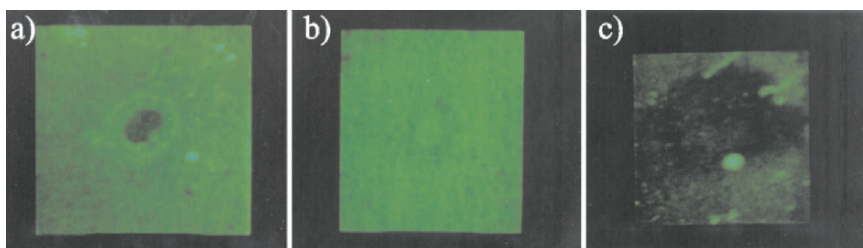


Fig. 8.11. Image of a dent on the accumulating screen surface in secondary ions $^{23}\text{Na}^+$ at the initial time moment (a) and in $^{24}\text{Mg}^+$ (b) and $^{39}\text{K}^+$ (c) at the subsequent time moments in the process of layer-by-layer etching.

periphery. It is seen from Fig. 8.11 that the dent boundary is enriched by Na, whereas the content of Na at its center is considerably lower. At the initial time moment, there are no Mg and K in the dent.

In the course of the dispersion of a material, secondary ions Mg^+ begin to escape from the dent. In a certain time interval, the intensity of their flux from the dent becomes higher than that from the surrounding region (see Fig. 8.11, b). Potassium behaves itself analogously. The measurements of the isotope ratios of Mg revealed their noncorrespondence to the natural ones. Moreover, the isotope ratios of Mg turned out to be violated both in

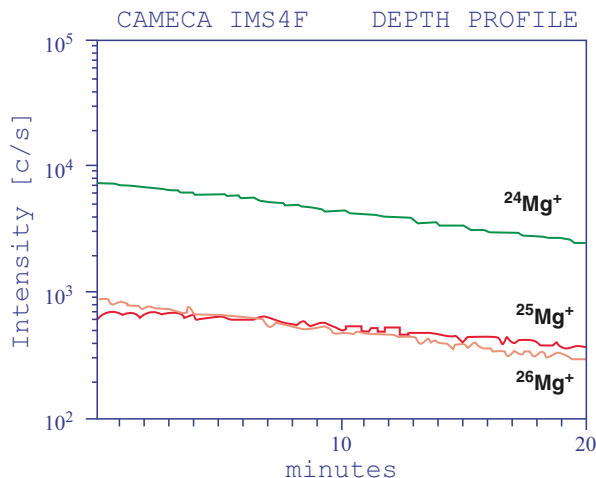


Fig. 8.12. Change of the distribution of the isotopes of Mg in a dent (see Fig. 8.11a, b) in the process of dispersion of its surface material.

the dent and in the surrounding region. The features of variations in the isotope composition of Mg during the dispersion of a material in the dent are shown in Fig. 8.12.

The coincidence of features of the behavior of Mg and K during the dispersion of the dent surface material allows us to assume that the origin of the latter is of the same nature as that of Mg. Therefore, we expected that the potassium would also reveal the anomalies of its isotope composition. However, it turned out that all K accompanying Mg during the dispersion of the upper layer of the dent surface has the natural isotope composition. In addition, not only potassium, but all other chemical elements (Si, Ca, S, and P) accompanying Mg had no anomalies of the isotope composition.

Then we performed the long-term etching of the dent surface, which completely removed the continuous layer containing Na, Mg, and K, and discovered a rounded small particle on its bottom enriched by K (see Fig. 8.11, c). The local analysis of this particle showed an anomalously low content of nuclide ^{41}K in it. The results of the determination of the distribution of nuclides ^{39}K and ^{41}K in it are given in Fig. 8.13, a. The ratio of nuclides $^{41}\text{K}/^{39}\text{K}$ registered in the particle varied from 2.7×10^{-3} to 3.5×10^{-3} (see curve 1, Fig. 8.13, b), which is lower more than by one order than the ratio for natural K which is equal to about 7.2×10^{-2} (see curve 2, Fig. 8.13, b). In Fig. 8.14, a, b, we present fragments of the high-resolution mass-spectra containing the mass-peaks of the ions of nuclides ^{39}K and ^{41}K . They illustrate the absence of the interference of their mass-peaks with the mass-peaks of complex and multiply charged ions.

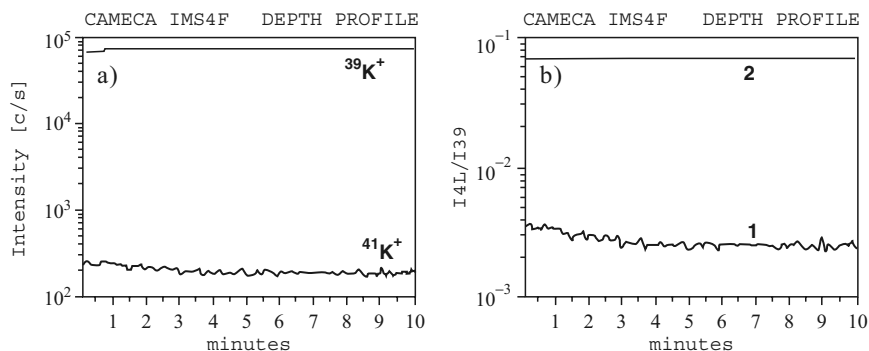


Fig. 8.13. Features of the variations in the distribution of the isotopes of K in a rounded particle (see Fig. 8.11, c) during its dispersion (a) and the comparison of the ratio $^{41}\text{K}/^{39}\text{K}$ registered in the particle (curve 1) with its value for natural K (curve 2).

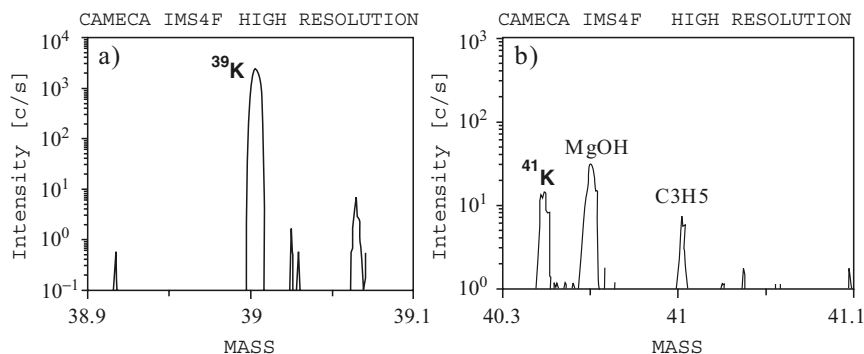


Fig. 8.14. Fragments of the mass-spectra of secondary ions registered in the mode with high resolution which contain the mass-peaks of nuclides ^{39}K (a) and ^{41}K (b), specimen No. 10623, the rounded particle.

8.2. Element Composition of Explosion Products

It is sufficiently easy to establish the fact that, due to the shock compression of a target, the products of its explosion that are observed on an accumulating screen contain chemical elements that were not present earlier in the initial materials of both the target and accumulating screen in registered amounts. For this purpose, we may take a target and an accumulating screen made of the same material. The last, if desired, should be quite pure (e.g., 99.99 mass. % Cu). In the study of the composition in the initial state by any local method (e.g., X-ray microanalysis), we can easily estimate which chemical elements, in what amounts, and in what form (inclusions, solid solution,

etc.) enter into it as admixtures. Then we can take the accumulating screen with target explosion products deposited on it after the experiment and study their composition by the same method. Despite the fact that, in this case, we added no new material to the accumulating screen (the target is manufactured of the same material), we found a significant difference of the composition of target explosion products from the composition of the initial copper. We discovered a huge number of particles, splashes, and films containing Na, Mg, Al, Si, P, S, Cl, K, Ca, Fe, Zn, Ag, Sn, Sb, Hf, Ta, W, Pt, Au, Hg, Pb, lanthanides, etc. Many of the indicated chemical elements in the initial material of the accumulating screen and target cannot be detected even upon a thorough search.

As for the problems concerning the determination of the most complete collection of “new” chemical elements appearing after the target explosion, the thorough study of their distribution over the accumulating screen surface and over depth, and the derivation of the exact evaluation of their content in the explosion products, their solution is considerably more complicated. To solve them, we combined the physical methods of determination of the element composition, including local and integral ones, together with the traditional methods of analytic chemistry.

8.2.1. Element Composition of Explosion Products by Physical Methods

The investigations concerning the study of the composition of target explosion products and the search of “new” chemical elements presumably appearing as a result of the explosion were carried on in the Laboratory, beginning from the very first stages of the Project. As one of the main objects of these investigations, we mention the layer of explosion products that deposits on an accumulating screen or remains in the target crater after the explosion. The main methods used for the solution of the indicated problems were, first of all, physical methods of determination of the composition, because they, as a rule, are more rapid and nondestructive as different from chemical methods of analysis of the composition. Among physical methods, we used X-ray electron probe microanalysis (XEPMA) (see Refs. 99, 105, 109–118), local Auger-electron spectroscopy (AES) (see Refs. 99, 102–104, 119–126), X-ray fluorescence analysis (XFA) (see Refs. 99, 110, 127–130), laser time-of-flight mass-spectrometry (LToFMS) (see Refs. 99, 105–108, 117), secondary-ion mass-spectrometry (SIMS) (see Refs. 99, 102–106, 117), and glow-discharge mass-spectrometry (GDMS) (see Refs. 99, 104, 105). From 1999 till 2004, we have applied the indicated methods to study more than 1100 specimens, on which we registered and calculated at least 16 000 spectra.

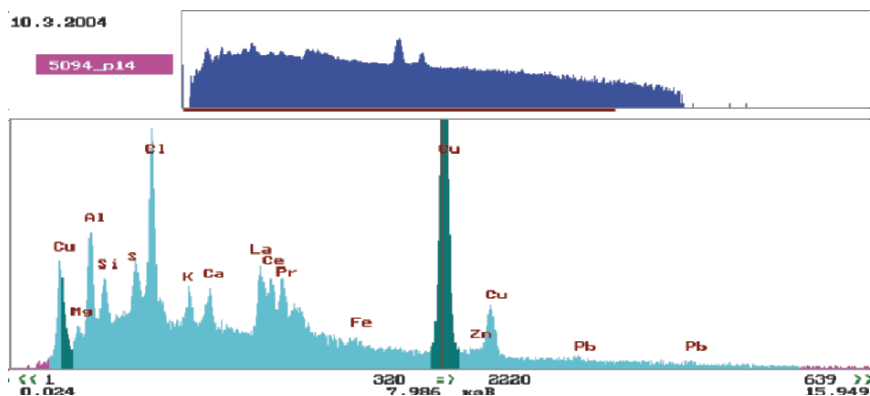


Fig. 8.15. X-ray spectrum registered by XEPMA on a globular inclusion contained in the explosion products of a copper target on a copper accumulating screen. Specimen No. 5094, the purity of the initial materials of the accumulating screen and target was 99.99 mass. %

In the study of the chemical composition of explosion products in every experiment (including those cases where the target and accumulating screen are manufactured of a single chemical element maximally purified from admixtures, for example, Cu (99.99 mass. %), Ag (99.99 mass. %), Pb (99.75 mass. %), and others), we registered up to several tens of chemical elements in considerable amounts that were not discovered earlier in the composition of the initial materials of targets and accumulating screens with the help of highly sensitive methods of investigation. Moreover, we found the chemical elements that were present in the form of admixtures, but in concentrations by three to seven and more orders of magnitude lower than those detected after the experiment (see Refs. 131, 132). To illustrate this fact, Fig. 8.15 shows the X-ray spectrum registered by XEPMA on a globular inclusion present on specimen No. 5094.

In the given spectrum, we see the characteristic X-ray peaks of 14 chemical elements. We note that the mentioned chemical elements include lanthanides. The last circumstance excludes practically the possibility for the indicated inclusion to be in the explosion products of the copper target as a result of contamination during the execution of the experiment, because lanthanides have very low abundance, only traces of them are present in common materials (see Ref. 133). It is worth mentioning that altogether 35 chemical elements not discovered in the initial materials of the target and accumulating screen were registered in the explosion products on specimen No. 5094.

According to the results of investigation by the method of scanning electron microscopy, the explosion products remained in the target crater

and deposited on the accumulating screens forming a layer consisting of irregularly distributed drops, splashes, films, particles, and other micro- and nano-objects with complex morphology (see Fig. 8.1 a, b). Because the explosion products are micro-objects, their study was first carried out mainly by local methods. At the beginning, we thought that local methods of analysis are most suitable for the study of such objects. Indeed, the studied chemical elements contained in micro-objects in slight amounts can be easily registered by local methods and, at the same time, be beyond the detection limits for even highly sensitive integral methods. However later, it became necessary to use also integral methods of analysis.

It turns out that the importance of the application of integral methods is especially high for the derivation of exact estimations of the amount of the substance appeared as a result of the target explosion. This is conditioned by the following. Only with the use of integral methods, by comparing the composition of the whole accumulating screen before and after the target explosion and taking into account the composition of a target material transferred on the accumulating screen, one can prove that the appearance of chemical elements in the explosion products, that were absent in the initial materials of the target and accumulating screen in considerable amounts, is a result of the nucleosynthesis, rather than a result of their redistribution from the volume of the accumulating screen or target. At the same time, we note that even if the use of integral methods simplifies, on the whole, the problem of evaluation of atoms-products of the nucleosynthesis, their results do not completely cover the results derived with local methods. On the one hand, the significant information given by local methods is the indication of the spreading area of appeared chemical elements and the character of their distribution in this area. On the other hand, the comparison of the estimations of a quantitative value, that are derived by different methods which complement one another, is very informative.

Local Methods of Determination of the Element Composition.

Amongst local methods, X-ray electron probe microanalysis and local Auger-electron spectroscopy were most frequently used for the determination of the composition of target explosion products. Both are nondestructive methods and allow one to determine the element composition of microvolumes of a substance rapidly and with high locality and accuracy. On the whole, these methods make it possible to analyze all chemical elements except for hydrogen and helium and cover the range of analyzed depths from about 2 nm to 3 μm (see Ref. 99). Among all problems of determination of the composition of target explosion products that were solved with the use of local methods, we separated two primary ones. The first problem consisted in the

determination, with the highest accuracy, the complete collection of chemical elements contained in the explosion products and in the estimation of the frequency of their appearance. In its solution, a special attention was paid to the determination of elements rarely occurring in the Nature. The main purpose of the second problem was the determination of the amounts of “new” atoms appeared in the explosion products as a result of the nucleosynthesis.

We show the result of solution of the first problem by the example of the methods of XEPMA and AES. By these methods alone, we investigated more than 950 specimens by the year 2004. The full collection of all chemical elements registered on the mentioned specimens and the frequency of their registration are given in Table 8.15. This table is shown in the form of the Periodic table of chemical elements, and its cells containing the discovered chemical elements in target explosion products are distinguished with gray color. The frequency of registration of an element is given in its cell under its designation. By analyzing the data in Table 8.15, we should like to note, first of all, the representative nature of the collection of chemical elements occurring in the explosion products. They fill practically completely the cells of the Periodic table of chemical elements. We also indicate the high frequency of the appearance of chemical elements with atomic numbers in the range from 22 to 30 in the explosion products. Their nuclei have high values of the specific binding energy per nucleon and almost maximum values of the specific binding energy per neutron (see Ref. 101). That is, they are nuclear structures that use neutrons in the collective process of nucleosynthesis in the most efficient way. By taking into account that the table includes only reliably registered elements, we call attention to the fact of the presence of almost all lanthanides, as well as Th and U, in its cells. In Fig. 8.16 one of these spectra containing characteristic X-ray peaks from L and M series of Th is presented.

Below, we illustrate the solution of the second problem by the example of applying the method of XEPMA to define the quantity of “new” atoms in the explosion products on the accumulating screen No. 5094. The X-ray microanalysis was performed on an X-ray microanalyzer REMMA-102 (SELMi, the town of Sumy, Ukraine) equipped with two wavelength-dispersion X-ray spectrometers and one energy-dispersion [with a Si(Li) detector] X-ray spectrometer. Spectra were registered at an accelerating voltage of the electron beam of 35 keV, the probe current of about 0.1 nA, and the residual pressure in the chamber with specimens of 2×10^{-4} Pa. The range of registered energies of the energy-dispersion spectrometer was 0.9 to 30 keV, the energy resolution on the line MnK_α at the counting rate of up to 1000 pulse/s was 150 eV, and the typical registration time of a spectrum was 200 to 400 s. For the quantitative analysis, we used a standard

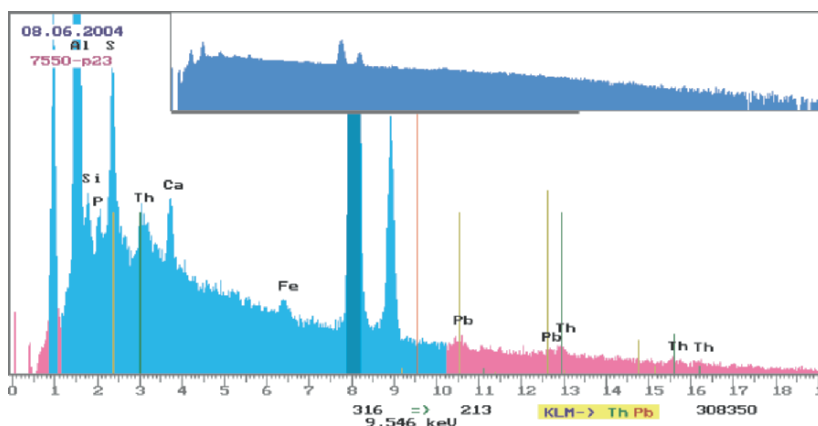


Fig. 8.16. X-ray spectrum registered by XEPMA on a globular inclusion contained in the explosion products of a copper target on a copper accumulating screen. Specimen No. 7550.

program supplied by the SELMI firm-producer of the device for the calculation of the concentrations of elements.

In the determination of the amount of atoms of the chemical elements appeared on an accumulating screen as a result of the target explosion by XEPMA, we took pure copper (99.99 mass.%) as the material of a screen and a target. We investigated exclusively the surface layer of accumulating screen. This accumulating screen was used as-received, i.e., it had not undergone any destructive cleaning or other procedures changing its composition prior to the analysis. The estimation of the amount of atoms of the “extraneous” chemical elements (all except for Cu) was carried out in two stages. First, we counted the number of atoms in the particles of explosion products lying on the surface screen and then in the enriched surface layer of the matrix of $2 \times 3 \mu\text{m}$ in thickness. The values derived in this manner were summed. The indicated value of the thickness of the enriched surface layer of the matrix followed from the data on the element concentration distribution over the depth of accumulating screens derived by SIMS.

We now describe the procedure and results of the first stage. The scheme of analysis on the first stage is given in Fig. 8.17. The analyzed area was a raster (a square) $54.3 \times 54.3 \mu\text{m}$ in size. We counted the number of all the present particles on it, analyzed their composition with an acute probe, and determined the number of atoms for each “extraneous” chemical element in each indicated particle. Then the analyzed area was moved by a distance equal to its side along the lines of analysis that formed an angle of 60° between themselves (see Fig. 8.17). In the realization of this procedure, we registered 417 spectra from different particles. By virtue of axial symmetry,

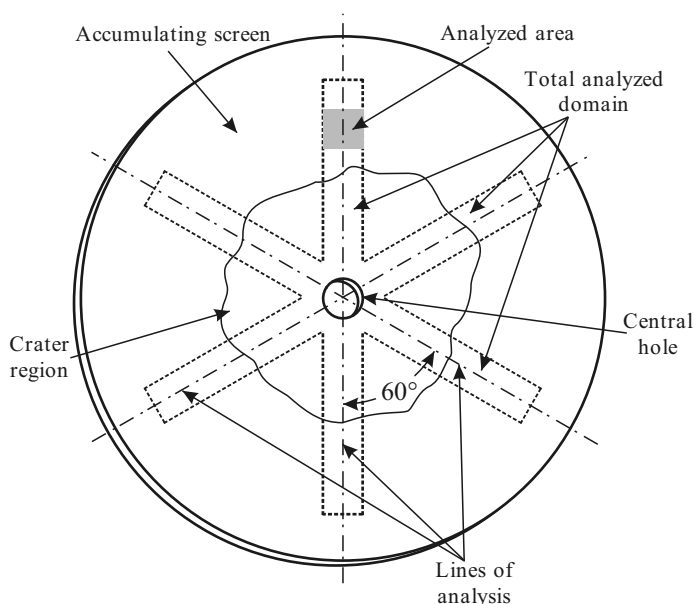


Fig. 8.17. The scheme of analysis upon the calculation of the total number of particles located on the surface of the entire accumulating screen No. 5094.

we assumed that all the lines of analysis are equivalent and the area analyzed along them is representative for the entire accumulating screen.

The total number of particles located on the surface of the entire accumulating screen, $N_{p.s.}$, was estimated as

$$N_{p.s.} = \frac{S_s}{S_a} \times N_{p.a.} \approx 2.0 \times 10^5, \quad (8.2)$$

where S_s and S_a are, respectively, the total areas of the screen surface and analyzed domain of the screen; $N_{p.a.} = 417$ is the total number of analyzed particles. Using the numbers of atoms N_{ij} of the i -th “extraneous” chemical element in the j -th particle found in the process of measurements, we calculated the amount of the i th “extraneous” chemical element in a particle of the averaged composition by the following formula:

$$\bar{N}_i = \frac{\sum_{j=1}^{N_{p.a.}} N_{ij}}{N_{p.a.}}. \quad (8.3)$$

With regard to Eq. 8.3, it is easy to get the estimation of the amount of atoms of the i -th “extraneous” chemical element contained in the particles located on the whole area of the accumulating screen surface. It is obvious that, for the indicated amount, the relation

$$N_i = \bar{N}_i \times N_{p.s.} \quad (8.4)$$

Table 8.16. Numbers of atoms of the “extraneous” chemical elements located in the particles positioned on the entire surface of accumulating screen No. 5094.

Element	Number of atoms	Element	Number of atoms
Mg	$3.06 \cdot 10^{15}$	Y	$2.04 \cdot 10^{14}$
Al	$9.08 \cdot 10^{16}$	Zr	$2.75 \cdot 10^{13}$
Si	$3.19 \cdot 10^{16}$	Ag	$6.14 \cdot 10^{15}$
P	$9.07 \cdot 10^{15}$	Cd	$2.20 \cdot 10^{15}$
S	$1.94 \cdot 10^{16}$	In	$1.92 \cdot 10^{15}$
Cl	$6.70 \cdot 10^{16}$	Sn	$1.61 \cdot 10^{16}$
K	$2.19 \cdot 10^{16}$	Te	$1.39 \cdot 10^{15}$
Ca	$1.28 \cdot 10^{16}$	Ba	$2.43 \cdot 10^{15}$
Ti	$3.48 \cdot 10^{15}$	La	$7.16 \cdot 10^{14}$
V	$5.08 \cdot 10^{13}$	Ce	$2.51 \cdot 10^{15}$
Cr	$2.40 \cdot 10^{15}$	Pr	$1.52 \cdot 10^{14}$
Mn	$5.89 \cdot 10^{14}$	Ta	$4.15 \cdot 10^{15}$
Fe	$5.11 \cdot 10^{16}$	W	$2.27 \cdot 10^{16}$
Co	$3.88 \cdot 10^{14}$	Au	$5.67 \cdot 10^{15}$
Ni	$2.07 \cdot 10^{14}$	Pb	$1.90 \cdot 10^{17}$
Zn	$2.87 \cdot 10^{16}$	Total:	$5.99 \cdot 10^{17}$

holds. The values derived as a result of the processing of spectra according to the above-mentioned procedure are given in Table 8.16. Finally, by adding all the values in Table 8.16, we get the total number of atoms of all the “extraneous” chemical elements contained in particles that are located on the accumulating screen surface as

$$N_{\Sigma p.} = \sum_i N_i \approx 5.99 \times 10^{17}. \quad (8.5)$$

By an analogous scheme, we determined the amount of atoms of the “extraneous” chemical elements contained in the enriched surface layer of the matrix of an accumulating screen. In this case, the raster was $11 \times 11 \mu\text{m}$ in size. The lesser area of the raster was chosen in order that we can easier search for parts of the surface screen without particles along the lines of analysis (see Fig. 8.17). The registration of spectra was realized from the whole raster area, and the number of these spectra registered from different

Table 8.17. Numbers of atoms of the “extraneous” chemical elements contained in the surface layer of the matrix of whole accumulating screen No. 5094.

Element	Number of atoms	Element	Number of atoms
Al	$2.13 \cdot 10^{17}$	Ca	$5.44 \cdot 10^{15}$
Si	$6.62 \cdot 10^{16}$	Mn	$8.77 \cdot 10^{14}$
P	$1.09 \cdot 10^{16}$	Fe	$3.57 \cdot 10^{16}$
S	$1.01 \cdot 10^{17}$	Zn	$4.63 \cdot 10^{16}$
Cl	$7.60 \cdot 10^{16}$	Pb	$6.22 \cdot 10^{15}$
K	$3.10 \cdot 10^{16}$	Total:	$5.93 \cdot 10^{17}$

areas was finally equal to 113. In the used scheme of analysis, the elementary analyzed domain was the surface layer of an accumulating screen $11 \times 11 \mu\text{m}$ in area (the size of a raster) and about $3 \mu\text{m}$ in thickness (the output depth of the registered X-ray emission of a specimen (see Ref. 99)). As above, we calculated first, by the results of separate measurements, the amount of the i -th “extraneous” chemical element in the elementary analyzed domain with averaged composition. Then, by knowing the ratio of the area of the elementary analyzed domain and the screen surface, it was recalculated into the amount of the i -th “extraneous” chemical element contained in the surface layer of the matrix of the whole screen. The results obtained in such a way are presented in Table 8.17.

As above, having summarized the values from Table 8.17, we get the total amount of atoms of the “extraneous” chemical elements contained in the surface layer of the matrix of the whole accumulating screen as

$$N_{\Sigma m.} \approx 5.93 \times 10^{17}. \quad (8.6)$$

Thus, with regard to Eqs. 8.5 and 8.6, the total amount of atoms of the “extraneous” chemical elements appeared on accumulating screen No. 5094 as a result of the target explosion and the number of nucleons contained in them are as follows:

$$N_{\Sigma} = N_{\Sigma p.} + N_{\Sigma m.} \approx 1.2 \times 10^{18}, \quad (8.7)$$

$$N_{\text{nucl.}} = 8.33 \times 10^{19}. \quad (8.8)$$

Analyzing the results obtained (see Eqs. 8.5 and 8.6), we note that the character of the distribution of “extraneous” atoms appeared on the surface of an accumulating screen is such that exactly one half of them is contained

in particles lying on the surface. The other half belongs to the surface layer of the matrix of at most $3\text{ }\mu\text{m}$ in thickness. It is easy to calculate that the mentioned amount of “extraneous” atoms corresponds to their concentration in the analyzed surface layer of the accumulating screen of about 3 mass. %. With regard to the purity of the initial materials of the accumulating screen and target (99.99 mass. %), this means that the obtained value must not be corrected onto the content of an admixture in both the initial material of the screen and the transferred material of the target.

Thus, the obtained value does correspond to the number of “extraneous” atoms appeared on the accumulating screen as a result of the target explosion. However, if we conclude in view of its value that the appeared “extraneous” atoms can be created only as a result of the nucleosynthesis, such a thought can be hardly classified as quite strict. In our opinion, this “logic” has two relatively weak points. On the one hand, the obtained value is not the result of a direct measurement, but is based on a number of statistical hypotheses and model ideas of the morphology of the surface layer of an accumulating screen, whose high degree of reliability and adequacy can be generally called under question. In other words, it is very difficult to estimate the accuracy of the obtained value. However we do believe the order of magnitude was estimated correctly, without little doubt. The last assertion follows from the fact that the estimation derived here agrees well with the results of solving the same problem by integral methods (see below), within which the required value is determined by means of direct measurements. On the other hand, the very procedure used here for the derivation of the required value does not theoretically exclude the possibility for “extraneous” atoms to appear on the surface of an accumulating screen from the bulk of the very accumulating screen by a redistribution of admixtures as a result of the explosion action, rather than due to the nuclear regeneration of the target material. In this connection, we note that it will be shown below by using integral methods that the evidence of the effect of redistribution of admixtures in the accumulating screen volume upon the target explosion is absent.

Integral Methods of Determination of the Element Composition.

The main goal of the application of integral methods is the derivation of a correct exact estimation of the number of atoms appeared in target explosion products as a result of the nucleosynthesis. To analyze the element composition, we used a highly sensitive glow-discharge mass-spectrometer (argon plasma) VG 9000 (VG Elemental, UK). The current and voltage of discharges were 1.8 mA and 1.1 kV, respectively. As a holder, we took a cell for plane specimens without cooling which ensured the diameter of an

analyzed domain to be 5 mm. The residual pressure in the chamber with specimens and in the spectrometer was at most 1×10^{-2} Pa and 1×10^{-5} Pa, respectively. For the ion beam, we used an accelerating voltage of 8 kV. The spectrometer has the range of analyzed masses 1 to 250. Its mass resolution $M/\Delta M$ is at the level of 7000 to 9000 on half the height of the mass-peak of Cu. For the quantitative analysis, we used a standard program for the calculation of the concentrations of elements supplied by the VG Elemental firm-producer of the device.

Here, we try to take into account all the above-mentioned drawbacks of the application of local methods for the solution of this problem. The method of GDMS by its essence is an integral one (the diameter of an analyzed domain was chosen to be 5 mm), and its application to derive the number of “extraneous” atoms appeared on accumulating screens requires no use of any model ideas of the morphology and structure of their surfaces. In other words, in the determination of a required value, the method can be used so that this value can be the result of a direct measurement. As for the effect of redistribution of the composition of an accumulating screen, it can be taken into account if, for example, the scheme of analysis is constructed so that, as a result of its application, we can register the composition of the whole accumulating screen, rather than the composition of a surface layer. In this case, the accumulating screen composition should be registered twice: first, in the initial state, and then after the target explosion. If the indicated compositions do not differ (the case where the initial materials of the target and accumulating screen are identical), then the enrichment of the surface layer of a screen occurs due to the redistribution of its composition over the specimen volume. But if the content of minor elements in the composition of an accumulating screen increases after the target explosion, we may say about the appearance of the atoms of “extraneous” chemical elements in it and will try to estimate their amount.

Consider the proposed scheme of analysis. In Fig. 8.18, we schematically show the cross-section of an accumulating screen where the procedure of analysis begins from that side of the screen, on which the film of target explosion products is deposited. It is obvious that, depending on the time of etching, the depth of an analyzed domain h_a varies. For example, if the condition

$$h_a \leq h \quad (8.9)$$

holds, where h is the thickness of a film, the results of measurements reflect the film composition. In the case where h_a meets the condition

$$h < h_a < H, \quad (8.10)$$

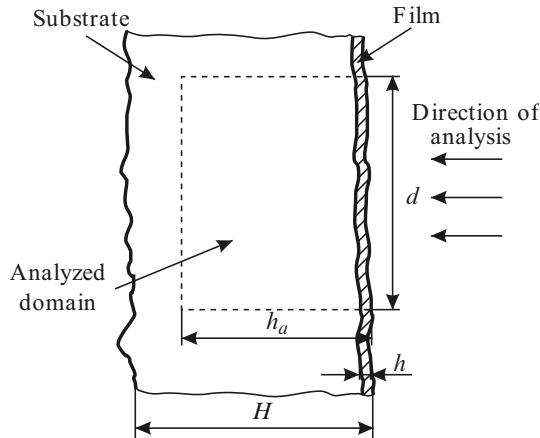


Fig. 8.18. Scheme of the cross-section of an accumulating screen.

(here, H is the thickness of a specimen) the film composition is dissolved more and more by that of the substrate material, and the results of measurements in this case have no significant physical sense. Finally, let us consider the situation where

$$h_a = H. \quad (8.11)$$

In this case, the results of measurements present the composition of an accumulating screen with a film. The case under study can be unambiguously characterized by the geometric factor k_0 which is the ratio of the area of a film cross-section S_f to the cross-section area of the whole accumulating screen S_0 . It is obvious that k_0 satisfies the relation

$$k_0 = \frac{S_f}{S_0} = \frac{d \times h}{d \times H} = \frac{h}{H}. \quad (8.12)$$

In correspondence with a real situation, we took the thicknesses of the film h and the specimen H to be equal to 2 and 500 μm , respectively. Then the geometric factor

$$k_0 = 0.004. \quad (8.13)$$

At first glance, it seems that we can determine the total composition of an accumulating screen following the above-presented scheme of analysis if relation Eq. 8.11 holds. However, its realization meets some problems. They are related to the fact that a mass-spectrometer with magnetic analyzer of masses is constructed in such a way that it can register, at every time moment, only those ions that are characterized by a specific nominal mass number defined by a value of the magnetic induction of the field in a mass-analyzer (see Ref. 99). Therefore, in order to register the whole mass

spectrum of a specimen, it is necessary to consequently scan the entire range of masses. In this case, in order that the registered spectrum have any physical sense, it is necessary that the specimen under study be homogeneous in its composition at least along the direction towards depth. Otherwise, we can fall in the situation where, for example, a specimen is etched through, but no mass-peak in the spectrum is registered. Such a situation could happen for a layered specimen, in which the layers consisting of pure chemical elements would be positioned from the specimen surface toward depth in the order of a decrease of their mass numbers, i.e., in the order inverse to one of the scanning of the mass range by a magnetic analyzer.

The above-mentioned difficulties can be avoided in the following way. We have manufactured the specimen as an assembly of several accumulating screens. The specimen was designed so that, on the one hand, it is homogeneous in composition in the direction from the analyzed surface towards depth, and, on the other hand, a domain analyzed on the mentioned surface has a geometric factor equal to k_0 . If two above-presented conditions are satisfied, the registration procedure of mass spectra is correct, and the results of analysis reflect the composition of an accumulating screen with a film. In Fig. 8.19, we schematically show a version of such an assembly of accumulating screens: a specimen of the "sandwich" type which is a stack composed from 20 to 30 accumulating screens closely adjacent one

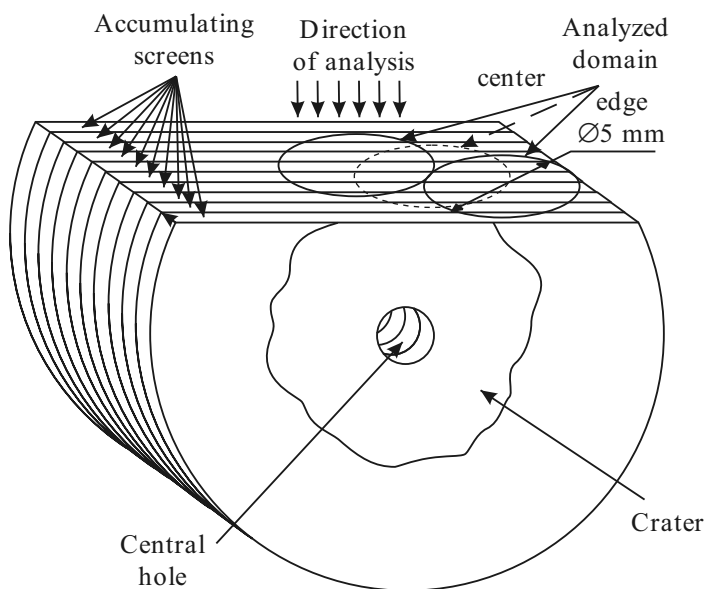


Fig. 8.19. Scheme of an assembly composed from accumulating screens which is a specimen of the "sandwich" type.

to another. In this case, the analysis was carried out from the lateral side of the assembly surface including the cut ends of accumulating screens (see “Direction of analysis” in Fig. 8.19). It is obvious that the proposed specimen of the “sandwich” type and the scheme of the execution of analyses on it satisfy the first above-formulated condition. Indeed, while moving from the analyzed surface towards depth even at a distance of the order of several hundreds of μm , we can consider the specimen to be homogeneous with a rather high accuracy. Moreover, the etching depth of a specimen upon the determination of its composition does not exceed $100\ \mu\text{m}$ in the real situation.

Finally, we need to clarify the situation with the geometric factor upon the determination of a composition on a specimen of the “sandwich” type. The corresponding scheme of analysis is presented in Fig. 8.20. It is obvious that the geometric factor of a “sandwich” satisfies the relation

$$k_s = S_{fs}/S_a, \quad (8.14)$$

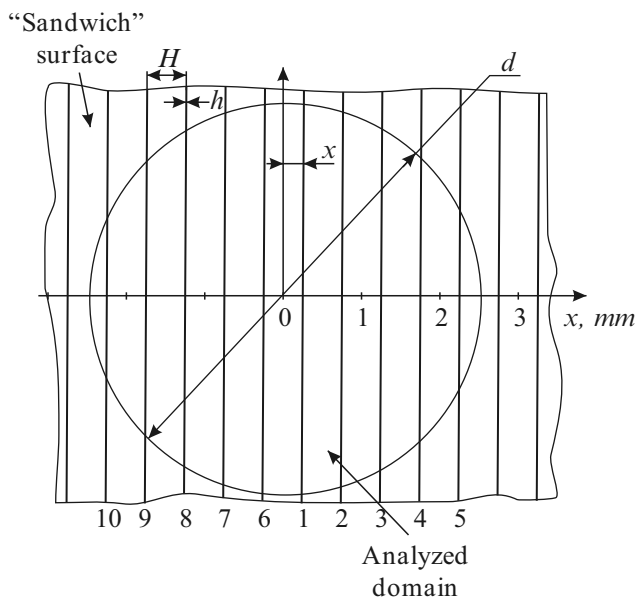


Fig. 8.20. Scheme of analysis on a specimen of the “sandwich” type, where H – the thickness of an accumulating screen, h – the thickness of a film from target explosion products, d – diameter of the analyzed domain, x – displacement of the film relative to the center of the analyzed domain, 1 to 10 – numbers of films got into the analyzed domain ($H = 500\ \mu\text{m}$, $h = 2\ \mu\text{m}$, $d \approx 5000\ \mu\text{m}$).

where S_{fs} is the cross-section area of all films located in the analyzed domain, and S_a is the area of the analyzed domain. Since $d = 5000 \mu\text{m}$, we get

$$S_a = \pi d^2/4 = 19634954 \mu\text{m}^2. \quad (8.15)$$

It is obvious that we can write

$$S_{fs} = hL, \quad (8.16)$$

where L is the length of all pellicular layers got into the analyzed domain. Now, the entire problem is reduced to the determination of L . In Fig. 8.20, all layers got into the analyzed domain are numbered from 1 to 10. We assume that the 1-st layer is displaced relative to the center of the analyzed domain by x . Then L can be written as

$$L(x) = \sum_{i=1}^{10} l_i(x), \quad (8.17)$$

where $l_i(x)$ is the length of the i -th layer which can be easily found for different i by using the Pythagorean theorem. Substituting Eq. 8.17 in Eq. 8.16, we can determine S_{fs} . It turns out that the quantity S_{fs} does not depend on x and, for any x from the interval $0 < x < H$, has the same value:

$$S_{fs} = 7.85 \times 10^4 \mu\text{m}^2. \quad (8.18)$$

Finally, substituting Eqs. 8.18 and 8.15 in Eq. 8.14, we find the value of the geometric factor k_s for the “sandwich”:

$$k_s = 0.004. \quad (8.19)$$

Comparing Eqs. 8.19 and 8.13, we see that

$$k_s = k_0. \quad (8.20)$$

Thus, we can conclude that the structure of specimens of the “sandwich” type and the proposed scheme of analysis on them satisfy two above-formulated conditions. This means that, with the help of a “sandwich”, we can correctly determine the general composition of an accumulating screen with a film.

We now describe the scheme and the procedure of calculation of the number of “extraneous” chemical elements appeared in target explosion products. Let C_{ti} be the concentration of the i -th chemical element in a target prior to its explosion, and let C_{si} be the concentration of the i -th chemical element of the initial material of an accumulating screen. These

values were derived by direct measurements before the experiment. The composition of the initial material of an accumulating screen was determined on the “sandwich” (see Fig. 8.19) collected from the blanks of accumulating screens undergone no high-energy actions. This allows us to reach a sufficiently high smoothing of inhomogeneities of a composition inherent in sheet materials (Refs. 134–136). It is worth noting that, in order to avoid the influence of surface contaminations on the results of analysis, the specimen surface was usually etched in the glow-discharge cell of a mass-spectrometer (argon plasma) for 20 to 30 min with a rate of $0.5\ \mu\text{m}/\text{min}$ in all measurements prior to the registration of a mass-spectrum. The usage of the mentioned procedure led to the removal of a surface layer of more than $10\ \mu\text{m}$ in thickness which contained usually an enhanced amount of admixtures. On the second stage of the measurement carried out on the “sandwich” collected from “processed” accumulating screens, we determined the general composition of the analyzed domain of an accumulating screen with target explosion products transferred on it. Let C_{sti}^* be the content of the i -th chemical element registered in the indicated case.

To calculate the concentrations of “extraneous” chemical elements appeared in the explosion products, it is convenient to consider that the experiment on the shock compression of targets is realized in two stages. We assume that the first stage after the target explosion involves only the transfer of its material on the accumulating screen and the target material preserves its initial composition in this case (see Fig. 8.21, a). The second

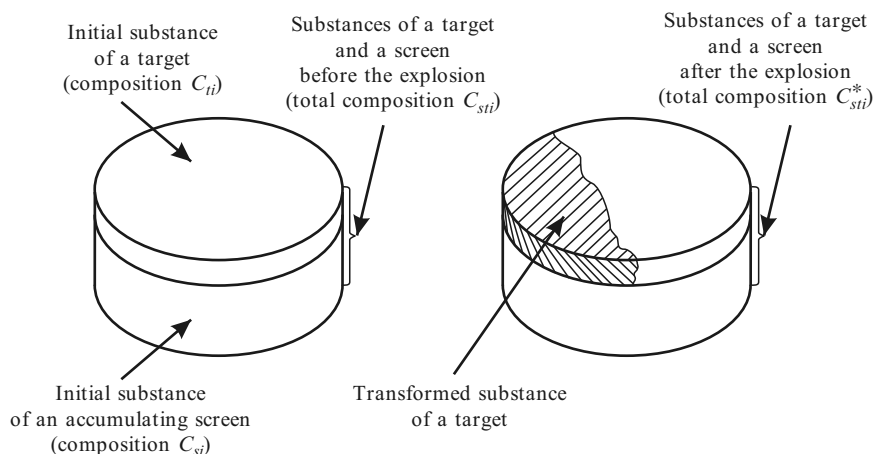


Fig. 8.21. Scheme of the analyzed domain of an accumulating screen with target explosion products transferred onto it [prior to process of nucleosynthesis (a) and after it (b)] clarifying the calculational procedure for the amount of appeared “extraneous” chemical elements.

stage is characterized by the processes of nuclear transmutation running in the target material transferred on the accumulating screen which lead to the appearance of “extraneous” chemical elements in it (see Fig. 8.21, b). Then, to describe the general compositions of the transferred target material and an analyzed domain of the accumulating screen at the end of the first stage (Fig. 8.21, a), it is convenient to introduce concentrations C_{sti} . It is obvious that the concentrations of “extraneous” chemical elements appeared due to a nuclear transmutation of the target material satisfy the following relation in terms of the above-introduced quantities:

$$C_{fi} = C_{sti}^* - C_{sti}. \quad (8.21)$$

Here, the symbol i is not referred to the major chemical element of the initial materials of a target and an accumulating screen. We recall that the quantities C_{sti}^* in Eq. 8.21 are known, because they are measured on “sandwiches” collected from “processed” accumulating screens. Thus, it is seen from Eq. 8.21 that, in order to find the concentrations of “extraneous” chemical elements, we must search for only concentrations C_{sti} . It is obvious that, in the case where the initial materials of a target and an accumulating screen are identical, we have the relation

$$C_{sti} = C_{si} = C_{ti}, \quad (8.22)$$

which allows us to solve Eq. 8.21 with respect to the unknown quantities C_{fi} .

We now consider the case where the initial materials of a target and an accumulating screen are different. In this case, the concentrations C_{sti} satisfy the relation

$$C_{sti} = \alpha C_{ti} + (1 - \alpha) C_{si}, \quad (8.23)$$

according to the mixture rule, where α is a share of the transferred target material in the total mass of the complex consisting of the transferred material of the target and that of the accumulating screen (see Fig. 8.21, a). Thus, relation Eq. 8.23 reduces the problem of determination of the concentrations C_{sti} to that for the coefficient α . To find the coefficient α , we will analyze the change in the composition of the complex shown in Fig. 8.21 under the transition from the first stage (a) to the second one (b). That is, we will try to find the interconnection of the concentrations C_{sti} and C_{sti}^* . As for the concentrations of minor chemical elements of a target and a screen, we may only conclude that they should not, at least, decrease, because their amount can only grow at the expense of products of the nucleosynthesis. As for concentrations of the major chemical elements of a target and a screen (let, for the sake of definiteness, the major elements of a target and an accumulating screen be, respectively, lead and copper), it will decrease in the

first case and remain constant in the second one. Indeed, the concentration of lead must decrease, since it is the major chemical element undergone a nuclear transmutation. At the same time, the concentration of copper does not change, because copper does not participate in nuclear reactions, being fully present in the accumulating screen (see Fig. 8.21). Hence, for the major chemical element of an accumulating screen, we can write

$$C_{stCu} = C_{stCu}^* \quad (8.24)$$

By using Eq. 8.24, we can now solve Eq. 8.23 at

$$i \Leftrightarrow Cu \quad (8.25)$$

and find the unknown coefficient α . Having known the coefficient α , we can now easily determine all the collection of C_{ist} . Substituting the derived values of C_{ist} in Eq. 8.21, we get the concentrations of all the “extraneous” chemical elements which have appeared in the target explosion products due to the nucleosynthesis.

By analyzing the question about both the scheme of measuring experiment and the procedure of calculation of the concentrations of “extraneous” chemical elements, we indicate one more circumstance, namely, both are developed in such a way that they exclude both the contribution of a number of sources of contamination to the required quantity and the influence of certain effects on it. In particular, we note once again that even if we are faced with the effect of an arbitrary redistribution of the admixtures earlier contained in the accumulating screen as a result of the target explosion (e.g., all they go onto the accumulating screen surface), this effect cannot change the amount of derived “extraneous” chemical elements. This circumstance is due to the fact that, in the used procedure, we subtracted the amount of all admixtures contained in the initial accumulating screen (see Eqs. 8.23, 8.21) from the amount of all minor chemical elements contained in the complex formed by target explosion products and the material of an accumulating screen (see Fig. 8.21). We note that, in this case, it is of no importance, where these admixtures were exactly in the initial accumulating screen and where they went after the target explosion. An analogous assertion is also true relative to the admixtures contained in the initial material of the target and to the products of its explosion which were transferred on the accumulating screen. The use of Eqs. 8.23 and 8.21 excludes them from the derived “extraneous” chemical elements. In other words, all the minor chemical elements of the initial materials of a target and an accumulating screen do not enter the determined amount of “extraneous” chemical elements.

Moreover, the proposed procedure also excludes the contribution of a number of unavoidable contaminations to the required value. As contaminations, we mean any entrances of an extraneous substance from the outside to

the “initial material of an accumulating screen–initial material of a target” system. It is obvious that the unavoidable entrance of an extraneous substance to this system occurs, for example, on the exposure of accumulating screens in air by means of adsorption or during the process of manufacturing of “sandwiches” from accumulating screens. In the latter case, we took precautions in order that the contributions from some sources of contamination be excluded completely and those from others be extremely minimized. The main point consists in that we manufactured “sandwiches” always under the same conditions. The exclusion of the contributions from contaminations of the indicated type to the derived amount of “extraneous” chemical elements happens for the reason that they are identical for the initial and processed “sandwiches” and, hence, are subtracted in calculations (see Eqs. 8.21, 8.23). In other words, the same contributions from contaminations to the initial and processed “sandwiches” behave themselves as if they were admixtures of the initial material of accumulating screens. Taking the above into account, it becomes clear that the amount of “extraneous” chemical elements derived according to the proposed method is affected by the contributions of only those contaminations which are present on the processed “sandwiches” and, at the same time, are absent on the initial “sandwiches”. In experiments, we met such a source of contamination, namely, the condensation of a substance on the accumulating screen surface upon the target explosion from the residual atmosphere of a vacuum chamber of the experimental setup. It is easy to estimate the order of magnitude of this source. In the residual atmosphere of a vacuum chamber of the experimental setup with a volume of 0.7 dm^3 and at a residual pressure of about 10^{-3} Pa , there is at most $1 \cdot 10^{-5} \text{ mg}$ of substance (vapors of a working liquid, rarefied air, hydrocarbons, etc.). Discussing the sources of contamination, it is necessary to note the fact that no transfer of a substance from the body walls of the experimental chamber on the accumulating screen occurs in the performed experiments, since we took certain technical measures to suppress the ricochet of the dispersed substance of a target from the body walls to the screen. The efficiency of the suppression of the mentioned process is evidenced by, both, the character of the directivity of splashes to the accumulating screen (see Fig. 8.1, b) and the absence of any significant correlation between the composition of explosion products deposited on the accumulating screen and the composition of a material of the chamber walls.

Finally, it is worth making some comments as for the accuracy of the determination of the amount of “extraneous” chemical elements. In calculations, for all the experiments, we used the values of C_{ti} , C_{si} , and C_{sti}^* averaged over the results of four measurements. Based on the results of repeated measurements, we may conclude that the total amount of “extraneous”

chemical elements was determined in each experiment with accuracy of at least 10 to 15 relat. %. At the same time, it is necessary to note that the accuracy of the determination of the amount of separate “extraneous” chemical elements, especially those with a low content in the nucleosynthesis products, could be considerably less. However, due to a smallness of the very value, such chemical elements gave no significant contribution to the error of the determination of the total amount of the atoms of “extraneous” chemical elements. Thus, we may conclude that, according to the given method of determination of the amount of products of the nucleosynthesis which are generated upon the explosion of a target and transferred onto the analyzed domain of an accumulating screen, it is necessary to subtract the amount of a condensate from the residual atmosphere of a vacuum chamber from the derived amount of all “extraneous” chemical elements. The derived value can be considered significant if it exceeds 10% to 15% of the amount of “extraneous” chemical elements.

In the present work, we describe the results of determination of the amount of products of the nucleosynthesis by the proposed method for four different experiments. For each of the experiments, the data on the main characteristics of the used “sandwiches” and targets are given in Table 8.18. Below, we make some comments on the presented data. As seen from Table 8.18, in the first experiment, the processed “sandwich” was collected from accumulating screens of 500 μm in thickness, on one side of which the explosion products of targets were positioned as layers of about 2 μm in thickness. In this case, the “sandwich” has a rather low value of the geometric factor, namely, 0.004, and the domain analyzed with a mass-spectrometer

Table 8.18. Main characteristics of the “sandwiches” and targets under study in experiments on the determination of the amount of products of the nucleosynthesis.

No.	“Sandwich”					Target	
	Mat. ^a	Purity, mass. %	Thickness of a screen, μm	Layers ^b	Geometric factor	Mat. ^a	Purity, mass. %
1	Cu	99.96	500	1	0.004	Pb	99.96
2	Cu	99.63	200	2	0.02	Ag	99.95
3	Cu	99.63	200	2	0.02	Al	99.93
4	Nb	99.80	100	2	0.04	Fe	99.90

^amaterial the detail was made of
^bnumber of layers with explosion products per screen.

covers only 10 ends of accumulating screens (see Figs. 8.19, 8.20). After the execution of the first experiment, it becomes clear that a low value of the geometric factor leads to that we register a rather low value, namely, a weak signal from the layer of explosion products of about $2\text{ }\mu\text{m}$ in thickness against the background of a strong signal from the accumulating screen of $500\text{ }\mu\text{m}$ in thickness playing the role of a substrate. In other words, to increase the accuracy of the determination of the amount of “extraneous” chemical elements, it was necessary to search for possibilities for increasing the geometric factor of the used “sandwiches”. In the second and third experiments, the geometric factor of the “sandwich” under study was increased up to 0.02, i.e., by five times. We succeeded in reaching this limiting value of the geometric factor on copper accumulating screens by means of a decrease of the thickness of a screen H (see Figs. 8.18, 8.20) up to $200\text{ }\mu\text{m}$ (at lesser thicknesses, copper screens are destroyed upon the explosion of a target) and the application of two layers of explosion products with $h \approx 2\text{ }\mu\text{m}$ on it: one layer on each side. Thus, the domain analyzed with a mass-spectrometer covered already 25 ends of accumulating screens on the “sandwich” (see Figs. 8.19, 8.20) both for the initial and processed specimens. The used measures not only enhanced the level of a useful signal, but also improved the averaging of inhomogeneities of the composition of the initial sheet material of screens. Finally, in the fourth experiment, we managed to decrease the screen thickness up to $100\text{ }\mu\text{m}$ with the application of explosion products on both sides of the screen, which increases the geometric factor of a niobium “sandwich” to a level of 0.04. Moreover, the number of the ends of screens fallen in the analyzed domain of a mass-spectrometer reached 50.

Table 8.19 shows the data of the performed experiments on the amounts (in mass. %) of “extraneous” chemical elements and the initial substance of a target transferred as a result of the target explosion on the analyzed domain of an accumulating screen in cases where it is located at the screen center or on its edge. Upon the analysis of the presented data, first of all, we paid attention to the following fact. In the experiments on silver, aluminum, and iron targets, the content of target explosion products (both the regenerated and nonregenerated parts) transferred onto the analyzed domain of an accumulating screen was much more than that in the experiment with a lead target. We note that this fact, in the first turn, is conditioned by a considerably grown value of the geometric factor of “sandwiches” used in the latter experiments, rather than by the increase in the very transferred amount of target explosion products onto the analyzed domain of an accumulating screen in terms of mass. In other words, the increase in the geometric factor led to the significant decrease of a share of the mass of a substrate, whose role is played by an accumulating screen, in the total mass

Table 8.19. Data on the amount of “extraneous” chemical elements and the initial substance of a target transferred onto the analyzed domain of an accumulating screen as a result of the target explosion.

Material of target/sandwich	Pb/Cu	Ag/Cu	Al/Cu	Fe/Nb
Content of “extraneous” chemical elements transferred onto the central domain of an accumulating screen, %	0.10 ± 0.02	0.60 ± 0.03	0.22 ± 0.02	1.35 ± 0.05
Amount of the initial substance of a target transferred onto the central domain of an accumulating screen, %	0.69 ± 0.03	5.36 ± 0.09	1.26 ± 0.05	3.74 ± 0.07
Total amount of the substance of a target transferred onto the central domain of an accumulating screen, %	0.79 ± 0.04	5.96 ± 0.09	1.48 ± 0.05	5.09 ± 0.08
Ratio of the amount of the initial substance of a target at the central domain to that on the edge	9.8	4.4	1.6	3.5
Ratio of the amount of “extraneous” chemical elements at the central domain to that on the edge	9.8	4.4	1.6	3.5

of a domain analyzed by a mass-spectrometer (see Fig. 8.21). It is obvious that a growth of the content of target explosion products transferred onto the analyzed domain of an accumulating screen ensured an increase in the accuracy of its determination. We also note that if the geometric factor of the “sandwich” under study was not increased by five times in experiments on the aluminum target, the determination of the amount of “extraneous” chemical elements in this case would meet difficulties. As for the accuracy of the results presented here, we give one more comment. The values of errors for the measured quantities were derived with regard to the results of the repeated measurements. The values given in the first three rows of Table 8.19 were determined independently one from another. Therefore, despite the fact that, by physical sense, the third value is the sum of the two first ones, the error of its measurement does not equal the sum of their errors.

Remarkable are the last two rows in Table 8.19. Consider the first one. As seen, the amount of the initial substance of a target transferred onto the analyzed domain of an accumulating screen located at its center exceeds considerably the indicated amount in the case where the analyzed domain is positioned at the edge of an accumulating screen in all experiments (see Fig. 8.19). We add that, in all the cases, the domain of collapse of a target substance, which is a source of its entrance onto the accumulating screen, was located on the symmetry axis of an accumulating screen and was away from it at the same distance. In this case, the values of the solid angles, at which the analyzed domains at the center and at the edge of an accumulating screen were seen, differed by, at most, a factor of 1.5. It follows directly from the above that the spatial distribution of the dispersed substance of a target upon its explosion is not isotropic, namely, in the solid angle covered the central part of an accumulating screen, the so-called a crater domain (see Figs. 8.1, 8.17, 8.19), we observe the increased density of the flux of target explosion products. The crater looks like a pit on the surface of an accumulating screen of 10 to 15 μm in depth which is positioned at its central part and has a diameter of about 5 to 7 mm. Its formation is caused by the removal of the screen substance as a result of the target explosion.

There are also other facts supporting the above assertion. For example, let us assume that the density distribution of the initial substance of a target over the accumulating screen surface is homogeneous, and the value of surface density is equal to that registered at the center of an accumulating screen. Then we can easily calculate the amount of the initial substance of a target transferred onto the whole accumulating screen. Such calculations give the amount of the initial substance of a target on the whole accumulating screen that exceeds the total mass loss of a target at its explosion determined by direct weighting. The absurdity of the above conclusion yields that the assumption of homogeneity of the distribution of the initial substance of a target over the surface of an accumulating screen is wrong. In other words, we can surely assert that the character of distribution of explosion products over the accumulating screen surface is such as that shown by the continuous line in Fig. 8.22.

By analyzing the same row from Table 8.19, we may note that the ratio of the content of the initial substance of a target registered at the screen center to its content at the screen edge in the case of a lead target is greater than those for silver, aluminum, and iron targets. In this connection, the appropriate question arises: "Does the last fact mean that, in experiments on silver, aluminum, and iron targets, the dispersion of a substance after their explosion occurs more isotropically and, respectively, the explosion products are distributed more uniformly over the accumulating screen"? The

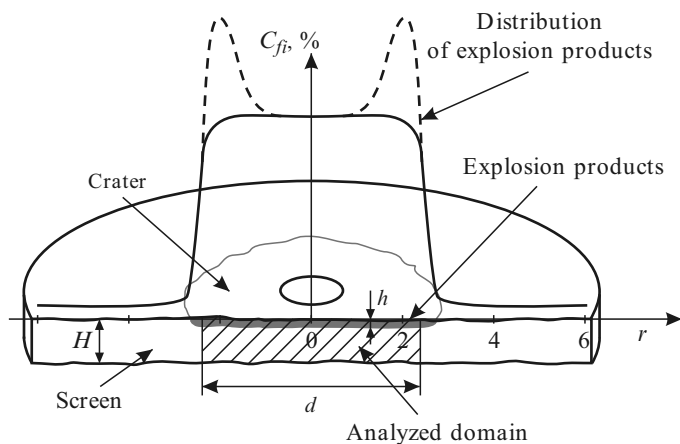


Fig. 8.22. Character of the distribution of “extraneous” chemical elements and the initial substance of a target over the surface of an accumulating screen. The dotted line shows the supposed enrichment at the crater edge.

answer to this question, despite the evident difference in the given values, is most likely negative. The reasons for the indicated differences consist probably in some nonequivalence of the used schemes of analysis in the cases under study that are demonstrated in Fig. 8.19. In experiments on lead targets, we managed to spatially separate the analyzed domains at the center of a “sandwich” and on its edge so that they do not practically overlap each other (see the continuous circles in Fig. 8.19). The analyzed domain on the edge of a screen does not also cover the crater region. Just this fact yields the large ratio of the content of the initial substance of the target at the center of the screen to its content on the edge. Due to thinner screens, “sandwiches” in experiments on silver, aluminum, and iron targets (see Table 8.18) have a considerably greater number of inner interfaces. The mentioned circumstance does not allow to reach the necessary hermeticity of the etching cell at a mass-spectrometer if the analyzed domain was located at the very edge of a “sandwich”. To get the necessary hermeticity, we displaced the analyzed domain somewhat to its center so that the domain partially overlapped the crater region (see the dotted circle in Fig. 8.19). This yields a greater value of the content of explosion products on the edge of a “sandwich” in the latter experiments.

On the other hand, it seems that the overlapping of the analyzed domains was not too large in the second case in order to decrease the ratio of the content of the initial substance of a target at the center of a screen to its content on the edge relative to the analogous ratio in experiments on lead targets by a factor of 3 to 5. The indicated circumstance is connected most likely to that the edge of the crater of an accumulating screen is more

enriched by explosion products as compared to their content at the central region of the crater (see the distribution shown by the dotted line in Fig. 8.22). Such a behavior of the distribution of the initial substance of a target over the surface of an accumulating screen allows us to easily explain the higher content derived on the edge of a “sandwich” in experiments on silver, aluminum, and iron targets. In other words, the behavior of the distribution of explosion products is most likely identical in both cases. At this point, we should like to make some comments. First of all, it seems to us that the distribution of explosion products fixed on an accumulating screen does not adequately reflect their distribution in the target substance running up to the accumulating screen. Their main difference consists in that a part of the substance of a target deposited on the central region goes away from its surface, which leads to the formation of a crater. Due to a lesser removal of the deposited mass from the edge of the crater of an accumulating screen, we derived the greater content of explosion products there as compared to that at the central region.

Consider now the last row in Table 8.19 and compare it with the previous one. In view of this comparison, we can state the complete coincidence of the characters of the distributions of the transferred initial substance of a target and “extraneous” chemical elements over an accumulating screen. This fact yields unambiguously not only the coincidence of the spatial positions of their sources, but the common mechanism of their transport from the mentioned sources onto an accumulating screen. Indeed, we exactly know how and whence the initial substance of a target appears on an accumulating screen and how this substance is distributed over it. Assume that the source of “extraneous” chemical elements registered on an accumulating screen is some source of contamination (in this case, it is of no importance whether a contamination arose during the target explosion in the experimental chamber of the setup, or in the process of preparation of specimens for the analysis). Then it is difficult to imagine that, each time for some reason, such a source “knew” the character of the distribution of explosion products of a target over an accumulating screen and developed in itself a mechanism of transport for its own contaminations onto a screen with a fine tuning to the prescribed character of their distribution. In other words, the considered flows of a substance have a unique source. It is obvious that such a source is a microvolume of the target substance undergone the shock compression, and “extraneous” chemical elements are none other than products of the laboratory nucleosynthesis.

The coincidence of the two last rows of Table 8.19 leads to the other extremely important conclusion, namely, that the ratio of the amounts of the initial substance of the target and “extraneous” chemical elements registered

at one analyzed domain of an accumulating screen is independent of the position of the indicated analyzed domain on the screen. The last conclusion can be somewhat generalized: in any solid angle with its tip at the point of collapse, the ratio of the amounts of the initial substance of a target and “extraneous” chemical elements upon the dispersion of explosion products is constant. This means that, upon the dispersion of the target substance, any separation of chemical elements is absent. The above assertion can be detailed. For example, let us compare the results of the determination of the amounts of separate “extraneous” chemical elements at the center and on the edge of a screen. In this case, we may note that their ratios are quite close to those of the total amounts of “extraneous” chemical elements or the initial substance of a target presented in Table 8.19. Moreover, they practically coincide with those for the “extraneous” chemical elements whose concentrations are great and are measured most exactly. The importance of the discovered absence of the separation of chemical elements upon the dispersion of the target substance consists in the following. This circumstance allows one to find the total amount of “extraneous” chemical elements or products of the nucleosynthesis originating upon the explosion of a target. Indeed, by direct measurements, we can determine a share of the amount of “extraneous” chemical elements in the total amount of target explosion products on an analyzed domain arbitrarily located on an accumulating screen. By virtue of the absence of the separation of chemical elements, the mentioned share of the total loss of the target mass determined by direct weighting is the total amount of products of the nucleosynthesis which originate during one explosion. We also note the circumstance that, however strange it seems, the total amount of products of the nucleosynthesis can be determined simply and with higher accuracy as compared to their total amount on the whole surface of an accumulating screen. To solve the last problem, it is necessary to know the exact form of the distribution of products of the nucleosynthesis over the surface of an accumulating screen, which is, in turn, a complicated problem.

Finally, we consider the parameters of the processes of transfer of a target substance and those of its nuclear regeneration in terms of masses, atoms, and nucleons. The corresponding data are given in Table 8.20. In its first part, we present the data on the diameter and mean mass loss of a target upon its explosion and on the masses of the analyzed domains of accumulating screens for various “sandwiches”. These are input data, and we will use them in the calculations of many parameters of the processes under study. Prior to the discussion of such a parameter as the target diameter, we consider its form. From the geometric viewpoint, a target consists of two parts: the cylindrical base with diameter D and the adjoining part that is

Table 8.20. Main parameters of the processes of mass transfer of a target substance and those of its nuclear regeneration.

	Material of target/sandwich	Pb/Cu	Ag/Cu	Al/Cu	Fe/Nb
I	Target diameter, mm	0.5	0.5	1.0	0.5
	Mean loss of the target mass upon the explosion, mg	2.60	3.10	5.30	1.66
	Mass of the analyzed domain of a screen, mg	87.38	34.95	34.95	14.84
II	Total mass of a target substance transferred onto the analyzed domain at the center of an accumulating screen, mg	0.690 ± 0.035	1.045 ± 0.017	0.259 ± 0.034	0.378 ± 0.016
	Share of the total mass of a target substance transferred onto the analyzed domain at the center of an accumulating screen in the target mass loss, %	26.5	33.7	4.9	22.8
	Mass of the regenerated substance of a target transferred onto the analyzed domain at the center of an accumulating screen, mg	0.090 ± 0.018	0.105 ± 0.005	0.038 ± 0.004	0.100 ± 0.004
	Share of the nucleosynthesis products in the target explosion products, %	13.0	10.0	14.5	26.5
	Number of regenerated atoms of a target transferred onto the analyzed domain at the center of an accumulating screen $\times 10^{18}$	1.66 ± 0.33	1.18 ± 0.09	0.388 ± 0.039	1.25 ± 0.05
	Number of nucleons contained in regenerated atoms transferred onto the analyzed domain at the center of an accumulating screen $\times 10^{19}$	5.50 ± 1.10	5.78 ± 0.29	2.26 ± 0.22	6.02 ± 0.22
III	Effective shortening of a target upon the explosion, mm	1.25	1.59	2.67	1.16
	Effective diameter of the nuclear regeneration zone of a target substance, mm	0.38	0.38	0.82	0.47

Table 8.20. *Continued.*

	Material of target/sandwich	Pb/Cu	Ag/Cu	Al/Cu	Fe/Nb
	Thickness of the shell of the initial material of a target surrounding the nuclear regeneration zone of a target substance, mm	0.06	0.06	0.09	0.02
	Total mass of the nucleosynthesis products of a target substance, mg	0.34 ± 0.07	0.31 ± 0.02	0.78 ± 0.08	0.44 ± 0.02
	Number of atoms contained in the total mass of the nucleosynthesis products of a target substance $\times 10^{18}$	6.27 ± 1.25	3.48 ± 0.17	7.96 ± 0.8	5.50 ± 0.20
	Number of nucleons contained in the total mass of the nucleosynthesis products of a target substance $\times 10^{20}$	2.07 ± 0.41	1.71 ± 0.09	4.64 ± 0.5	2.65 ± 0.10
IV	Mean number of nucleons per atom in the detected part of the regenerated substance of a target	33.1	49.0	58.3	48.2
	“Conditionally average product” of the detected part of the regenerated substance of a target	S	Ti	Ni	Ti

a hemisphere (see Fig. 8.23, “Target–anode”). As for the parameter of the mean mass loss of a target upon the explosion, we point out that, in order to manufacture “sandwiches”, we took usually from 20 to 70 accumulating screens depending on their thickness. For the majority of accumulating screens, we performed two shots at targets in order to apply explosion products on both sides of a screen. Therefore, the number of exploded targets was, as a rule, twice as many the number of screens used for the manufacture of “sandwiches”. Each time after the target explosion, we determined its mass loss by weighting. It is clear from the above that the mass losses of targets given in Table 8.20 are averaged over the results of many tens of measurements. The masses of analyzed domains were determined as shares of the mass of the whole accumulating screen which are proportional to

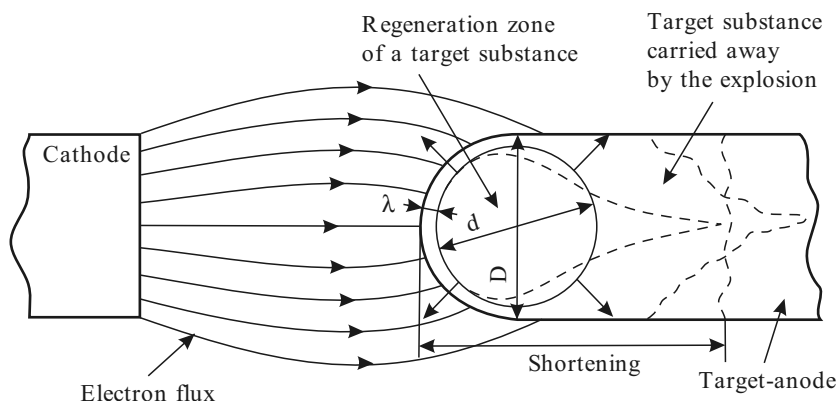


Fig. 8.23. Scheme of the formation of the superhigh-compression zone of a target substance clarifying the estimation of its size.

their volume shares (see Figs. 8.21, 8.22). We note that the mass of an analyzed domain is the basis for all values given in percentage and presented in Table 8.19.

Consider the second part of Table 8.20, where the target explosion is described with regard to the analyzed domain of an accumulating screen. We make some comments about how we derived the presented values. We calculated the total mass of a target substance transferred onto the analyzed domain at the center of an accumulating screen and the mass of its regenerated part as the corresponding percentage shares (see Table 8.19) of the mass of a substance contained in the volume of the analyzed domain of an accumulating screen. In those cases where we used accumulating screens containing target explosion products on both sides for the manufacture of “sandwiches”, we took only a half of the indicated percent shares. That is, all data given in Table 8.20 are calculated for one explosion of a target. The calculational procedure for values of the second row of this part of Table 8.20 is obvious in view of the physical sense of the value under question.

The fourth row was calculated with the use of the third and first rows of the same part of Table 8.20. In other words, it presents the share of nucleosynthesis products in target explosion products for the analyzed domain at the center of an accumulating screen. However, by virtue of the absence of the separation of chemical elements under the dispersion of a target substance, the value of the mentioned share can be referred not only to an arbitrarily positioned analyzed domain on the accumulating screen, but to the whole lost mass of a target upon its explosion. Therefore, the mentioned value has, indeed, sense of the share of nucleosynthesis products in target explosion products. As for the physical sense of the value under question, we note that it cannot be considered as the efficient or useful yield

of the process of nuclear regeneration of a target substance due to the action of a subrelativistic beam of electrons on it. Most likely, on the contrary, all the substance of the target domain, where the action of the incident beam of electrons is focused, undergoes a nuclear transformation (see “Regeneration zone of a target substance” in Fig. 8.23). Therefore, on an accumulating screen in explosion products, we register both the initial and regenerated substances of a target only for the reason that the very regeneration zone is positioned inside the head part of a target. The zone is surrounded by the shell of the initial material of a target with thickness λ (see Fig. 8.23), where the energy density does not reach the critical level and, hence, nuclear reactions do not run. However, upon the target explosion initiated inside the collapse zone, the content of the latter (products of the nucleosynthesis) is carried away from the target together with the content of the mentioned shell (see Fig. 8.23).

The adequacy of the presented ideas of the geometry of a target explosion and, in particular, of the position of the region of nucleosynthesis inside the volume of its head part is confirmed by the results of measurements of the isotope composition of target explosion products deposited on accumulating screens (see Sec. 8.1.2). For example, anomalies of the isotope composition of chemical elements are usually registered in the upper layer of explosion products. But, in the lower layer adjacent to an accumulating screen, we found more frequently the isotope composition corresponding to the natural abundance of chemical elements. This fact indicates without doubt that, prior to the dispersion of a target substance, the nonregenerated part of explosion products was external relative to nucleosynthesis products. That is, the zone of nucleosynthesis was surrounded by a shell of the initial material of a target. The last means that the zone was located in the inner region of the head part of a target.

It becomes clear from the above-considered elementary ideas of the process of explosion of a target substance that the share of nucleosynthesis products in target explosion products, η , is equal to the ratio of the total mass of nucleosynthesis products to the target mass loss, Δm . By assuming in the first approximation that the region with a regenerated substance of a target is a ball (see Fig. 8.23), we easily get that its diameter

$$d = \sqrt[3]{6\eta\Delta m/\pi\rho}, \quad (8.26)$$

where ρ — the density of the initial material of a target. The diameters of the zone of regeneration of a target substance calculated by Eq. 8.26 are given in the second row of the third part of Table 8.20 for experiments on various targets. Finally, the fifth and sixth rows of the second part of Table 8.20 present the numbers of regenerated atoms of a target transferred

onto the analyzed domain at the center of a screen and the numbers of nucleons contained in them. The mentioned values are calculated by using the concentrations of separate “extraneous” chemical elements transferred onto the analyzed domain at the center of an accumulating screen. However, these concentrations are not given here because of the huge volume of the data describing them.

In the third part of Table 8.20, we present the data characterizing the process of explosion of a target on the whole. In its first row, we demonstrate the data on the shortening of a target as a result of carrying away its substance upon the explosion. The shortening is easily calculated by using both the known geometry of a target in its initial state and the target mass loss. The essence of the second row of this part was described above. The third row of the third part of Table 8.20 shows the values of the parameter λ characterizing the thickness of the shell of the initial material of a target surrounding the zone of regeneration of a target material. This parameter can also be interpreted as the free path of a density wave generated by the electron beam incident on a target up to the time when the target substance, being on its leading edge, reaches the critical level of density sufficient for the running of nuclear reactions. The geometric reasons (see Fig. 8.23) yield that the mentioned parameter, under the assumption that the zone of nuclear regeneration of a target substance has the form of a ball, can be estimated by the formula:

$$\lambda = (D - d)/2, \quad (8.27)$$

where D is the the target diameter, and d is the diameter of the zone of regeneration of a target substance. The data in the fourth row with regard to the absence of the separation of chemical elements were calculated as the shares of the nucleosynthesis products (see the fourth row of the second part of Table 8.20) in the target mass loss. The data in the fifth and sixth rows were derived from those in the fifth and sixth rows of the second part of Table 8.20 by multiplying them by the ratio of the total amount of nucleosynthesis products to their amount transferred onto the analyzed domain at the center of an accumulating screen. Finally, by dividing the number of nucleons contained in the nucleosynthesis products by the number of atoms, in which they are located, we filled the last fourth part of Table 8.20.

Having described the procedure of derivation and having refined the physical sense of the quantities presented in Table 8.20, we move to the analysis of their values. As for the first part of Table 8.20, there is no large room for their analysis. In fact, they are the initial data and mostly turn out to be as-taken. Apparently, it is worth noting only the pronounced dependence of the target mass loss on the atomic number of the major chemical

element of its initial material. All looks as if the interconnection between the quantities under consideration were close to an inversely proportional one. However, we have no sufficient amount of data to reliably draw this conclusion. Most likely, the nature of the mentioned interconnection is much more complicated. It seems that the target mass loss should be affected by electric, thermophysical, elastoplastic, strength-related, and nuclear-physical properties of its initial material.

In Table 8.20, its second and third parts are principal. In the second part, the process of explosion of a target is characterized with regard to the analyzed domain at the center of an accumulating screen. This shows some boundedness of the data presented here. If the analyzed domain is positioned at any other place of an accumulating screen, we get entirely other values of the quantities under study, whereas the process of explosion of a target remains the same. That is, these data do not describe the process of explosion of a target on the whole. At the same time, it is worth noting that these data are most reliable, because they are the result of direct measurements and they should be compared with the power of the sources of contamination and with the results of determinations of the amount of nucleosynthesis products by local methods. Such a comparison showed that the amount of nucleosynthesis products registered only in the analyzed domain at the center of an accumulating screen exceeds the contribution of contaminations by about four orders, and their amounts derived by local and integral methods are in good agreement. The last means that we have measured reliably the amount of nucleosynthesis products, and the contribution of contaminations can be neglected.

By comparing the results of measurements in the experiments on various targets, we may state that the share of target explosion products transferred onto the central analyzed domain of an accumulating screen is about 20% to 30% in the majority of cases, which considerably exceeds the mentioned value for the isotropic distribution of the dispersed substance of a target. It follows from the above discussion that the main flow of a substance from the target is directed in the cases under study along the symmetry axis of an accumulating screen to its central region, to the crater. At the same time, the indicated share for an aluminum target is 4.9%, which corresponds to the distribution of target explosion products that is close to an isotropic one.

We can present several reasons for the absence of a pronounced directivity upon the dispersion of explosion products in the case of an aluminum target. The first one is illustrated by the scheme drawn in Fig. 8.24 with the observance of proportions. It shows a cross-section of the head part of a target by the plane passing through the center of a region, where the nuclear

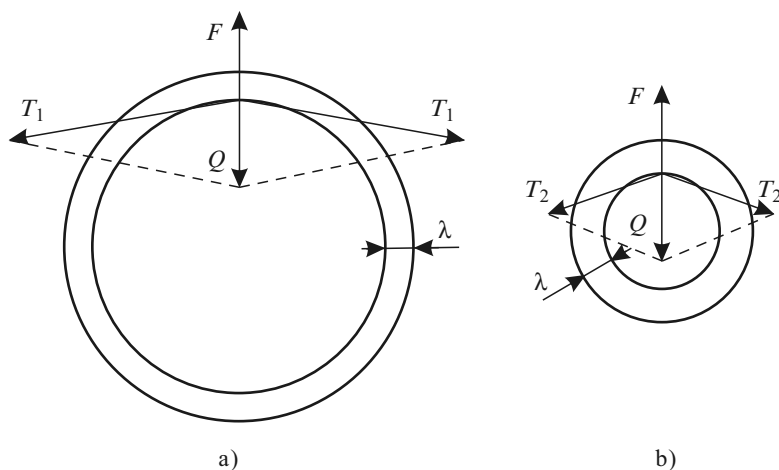


Fig. 8.24. Scheme of a cross-section of the head parts of targets by a plane passing through the center of a domain of nuclear regeneration of a target substance. An aluminum target – (a); lead, silver, and iron targets – (b).

regeneration of a target substance occurs (see Fig. 8.23). Aluminum targets correspond to Fig. 8.24, a, and lead, silver, and iron targets corresponds to Fig. 8.24, b. As seen from the presented scheme, the action of a force F outwards an aluminum target induces the appearance of tensile forces in every elementary volume of the target material and the relevant compensating forces T_1 . The latter in modulus considerably exceed analogous forces T_2 arisen in elementary volumes of the external shells of lead, silver, and iron targets due to the lesser curvature of the aluminum target shell. The appearance of tensile stresses in a shell can induce its early destruction, which will promote, obviously, the isotropic dispersion of explosion products. The second reason for the early destruction of an aluminum shell can consist in its low ability to the inertial confinement of the expanding material inside it. This reason is conditioned, first of all, by both a small value of the ratio of the thickness of the external shell to the diameter of the internal ball (see Fig. 8.24) and the low density of aluminum.

As for the share of nucleosynthesis products in target explosion products, it is equal to 10% to 15% in most cases and reaches the level of 26.5% only in experiments on iron targets. It is difficult to find the obvious reasons for this fact. At the same time, we note that the fact fits the following regularity: “The higher the specific binding energy (see Ref. 101) of nuclei of the initial material of a target, the greater is the share of nucleosynthesis products in its explosion products.” In the central analyzed domain of an accumulating screen, we registered about 40 to 100 μg of nucleosynthesis products which contained about 0.4×10^{18} to 1.7×10^{18} regenerated

atoms and 2.3×10^{19} to 6.0×10^{19} nucleons. We note a smaller amount of nucleosynthesis products registered in the central analyzed domain of an accumulating screen in the case of aluminum targets. This fact can be explained by that the lesser amount of nucleosynthesis products falls at the center of an accumulating screen due to the negligible directivity of the flow of target explosion products in the case under consideration, rather than by their reduced synthesis upon the target explosion (see the second row of the second part and the fourth row of the third part of Table 8.20).

The data presented in the third separated part of Table 8.20 characterize integrally the target explosion. We recall once more that most data are derived with regard to the assertion about the absence of a separation of chemical elements upon the dispersion of a target substance. This assertion is true quite exactly in all the cases under study (see two last rows in Table 8.19). By analyzing the data on the shortening of targets and on the diameters of the zone of nuclear regeneration of their substance, we pay attention, first of all, to the anomalously high values of the mentioned parameters for aluminum targets. These values are practically twice as much the corresponding parameters for lead, silver, and iron targets. It is obvious that this fact is conditioned by the double initial diameter of aluminum targets (see the first row of the first part of Table 8.20). Thus, it follows from the above that the shortening of a target and the diameter of the zone of nuclear regeneration of its substance depend linearly on the initial diameter of a target upon its variation, at least, in the interval from 0.5 to 1.0 mm. This yields that the masses of explosion products and nucleosynthesis products are proportional to the cube of the initial diameter of a target under its change in the same interval. We also note that, apparently, the essence of this dependence is exclusively defined by the geometry, because the nature of the initial material of a target has practically no considerable influence. The last circumstance is related, most likely, to that the thickness of the shell composed from the initial material of a target and surrounding the zone of regeneration of a target material, is less at least by one order in all cases than the target diameter (see the third row of the third part of Table 8.20). In other words, the mass of nucleosynthesis products appears as a cube on the diameter of a target by virtue of the fact that the conditions for the nuclear regeneration of a target material under the action of a subrelativistic beam of electrons on it appear in the close vicinity to its surface.

Let us compare these parameters derived in experiments on the same initial diameter of a target, i.e., let us omit the geometric parameter. In this case, it should be noted that the minimum shortening of a target and the maximum diameter of the zone of nuclear regeneration of its substance are

observed for iron targets. It is obvious that the values of these parameters for iron targets correlate with the already mentioned anomalously high share of nucleosynthesis products in their explosion products.

Finally, we make some comments in connection with the refinement of the physical content of the above-considered parameters. In Fig. 8.23, it is schematically shown that, after the carrying away of a target substance as a result of the explosion, the target has a more or less even end, and the target shortening consists in the reduction of its initial length by the length of its head part lost upon the explosion. Just this geometric sense is inherent in the parameters of shortening of a target given in the third part of Table 8.20. In reality, the end of a target after its explosion is not planar. In its central part, we usually observed a deep and narrow channel directed along the symmetry axis of the target (see the corresponding dotted lines in Fig. 8.23). The lateral walls of this channel are destroyed upon the explosion along the generatrices of the cylindrical surface of the target and have the form similar to that of the petals of a flower, by unbending from the cathode to the accumulating screen (see Fig. 8.23). As a result of the explosion, a real target acquires such an external form and structure as those of the target shown in Fig. 8.1, a. In other words, the mentioned parameter of the shortening of a target is an effective one, which is reflected in the corresponding part of Table 8.20.

Apparently, all the above-presented about the shortening of a target is also valid, to a great extent, for the zone of nuclear regeneration of a target substance. For simplicity, performing the calculations, we believe that the zone has the form of a ball (see the continuous line in Fig. 8.23). However, its form is similar, most likely, to that of a strongly elongated drop of liquid falling towards the cathode (see the dotted lines in Fig. 8.23). The assumption of such a form of the zone of nuclear regeneration of a target substance is supported by both the morphology of the fracture surface of the exploded target (the presence of a crater with deep channel, see Fig. 8.1a) and the results of experiments on the registration of the optical image and X-ray emission of a plasma bunch appearing at the place of the target explosion around its geometric center with the help of an obscure chamber and a fine-mesh collimator (see Chapter 6). Thus, like the previous case, we characterize the size of the zone of nuclear regeneration of a target substance by its efficient diameter (see Table 8.20).

Consider the last three rows of the third part of Table 8.20. It follows from the data given there that the total mass of nucleosynthesis products originating upon the explosion of a target substance is 310 to 780 μg , which corresponds to 3.5×10^{18} to 8.0×10^{18} regenerated atoms or 1.7×10^{20} to 4.6×10^{20} nucleons. We emphasize the results of experiments

on aluminum targets. There, a double amount of registered products of the nucleosynthesis is generated as compared to the experiments on lead, silver, and iron targets. We recall that the situation looked the opposite when we considered the analyzed domain at the center of an accumulating screen. Hence, to understand what occurs upon the explosion of targets, we must use the integral estimations of characteristics of the process.

Let us consider the high-level registered characteristics of the productivity of aluminum targets. We recall that their initial diameter was twice as much the diameter of other targets (see Table 8.20). If we take into account that the mass of nucleosynthesis products is proportional to the cube of the initial diameter of a target, we see nothing astonishing in the characteristics of productivity. Indeed, with the 0.5-mm initial diameter, we got only about $100\text{ }\mu\text{g}$ of nucleosynthesis products on aluminum targets, which is more than modest, relative to the results derived on other targets. In view of the comment made, we should set iron targets at the leading position by this parameter as well.

Finally, we will analyze the data of the last part of Table 8.20. Compare the results presented there for the experiments on the “lightest” and “heaviest” targets. It is striking that the effective representative of the nucleosynthesis products in the case of a lead target is sulfur, whereas for an aluminum target, is nickel. Moreover, the regularity “the heavier the atom of a target, the lighter is the atom of the effective representative of its nucleosynthesis products” is manifested quite clearly. This regularity contradicts the common ideas. It is obvious that a “heavy” element can be easier obtained due to the nucleosynthesis from a “heavy” target. Moreover, there is no complete clearness as for the inner nucleon balance. For example, the share of neutrons in the atom of an effective representative of nucleosynthesis products in experiments on aluminum targets is higher as compared to their share in atoms of the initial material of a target and is lower in the case of lead targets. Here, it is appropriate to remember the outstanding characteristics of iron targets in the generation of nucleosynthesis products and the minimum productivity of aluminum targets. We emphasize that this fact remains inexplicable in the framework of the existent ideas. Indeed, the nuclei of atoms of aluminum, possessing the huge stock of specific binding energy, could be “burned” in nuclear reactions up to nuclei Fe^{58} or Ni^{62} and give greater amounts of nucleosynthesis products and released energy. However, nothing is observed in experiments on aluminum targets. On the contrary, by virtue of the absence of a stock of specific binding energy in the nuclei of atoms of iron, it seems that nuclear reactions must not run on iron targets at all. But we observe the maximum amount of nucleosynthesis products with them.

Table 8.21. Estimations of the energy yield of nuclear transformations in experiments on lead, silver, aluminum, and iron targets.

Nuclear transformation	Specific binding energy*, MeV		Increment of specific binding energy, MeV	Number of nucleons	Total energy yield, kJ
	initial	final			
Pb $\rightarrow$ S	7.870	8.497	0.627	$2.07 \cdot 10^{20}$	$2.08 \cdot 10^4$
Ag $\rightarrow$ Ti	8.551	8.714	0.163	$1.71 \cdot 10^{20}$	$4.47 \cdot 10^3$
Al $\rightarrow$ Ni	8.332	8.748	0.416	$4.64 \cdot 10^{20}$	$3.09 \cdot 10^4$
Fe $\rightarrow$ Ti	8.787	8.714	-0.073	$2.65 \cdot 10^{20}$	$-3.10 \cdot 10^3$

* see Ref. 101.

We are also faced with paradoxes by estimating the energy yield of the nuclear transformations having occurred in the performed experiments. In the first three experiments, judging from the composition of nucleosynthesis products, we got a positive energy yield from reactions (see Table 8.21) with large values. The release of such an amount of energy cannot be overlooked. But the situation is the opposite for the last type of experiments: the energy yield is negative, which indicates the impossibility for the considered nuclear transformations to have occurred, whereas the experiment results demonstrate there must have been nuclear reactions.

The analyzed situation can be clarified only if we assume that we have not registered all products of the nucleosynthesis, but only their detectable and identified part: the chemical elements that belong to the known part of the Periodic table. It is not so obvious, but a conclusion we have been driven to by our data that the undetected part of nucleosynthesis products can be represented only by superheavy chemical elements. In the latter case, the dependence of the atomic number of an effective representative of nucleosynthesis products on the atomic number of the chemical element of the initial material of a target, variations in the inner nucleon balance in nuclear transformations, and extreme characteristics of the productivity of iron and aluminum targets are easily explained by a change in the relative mass shares of the detectable and undetected parts of nucleosynthesis products. We also indicate that the energy balance will be satisfied if the effective representative of the undetected part of nucleosynthesis products in experiments on iron targets has the specific binding energy per nucleus of at least

8.787 MeV and of at most 7.870, 8.551, and 8.332 MeV in experiments on lead, silver, and aluminum targets, respectively.

Thus, based on the results of the performed studies, we can conclude that an important problem is solved upon studying the composition of nucleosynthesis products within local and integral methods of analysis: we derive the quantitative characteristics of the detectable part of nucleosynthesis products. However, it turns out that this composes only a part of the problem of investigation of the composition of nucleosynthesis products. Its other part consists in the registration and study of the else undetected part of nucleosynthesis products. In other words, the performed studies have led to the search for superheavy chemical elements in the nucleosynthesis products which would seem to be the most dramatic findings of the project.

8.2.2. Element Composition of Explosion Products by Chemical Methods

Chemical methods of determination of the composition of objects are not universal and convenient, as the physical methods of analysis, and their share in the solution of analytic problems continuously decreases in recent years. Such a situation is conditioned, first of all, by the fact that the chemical methods, as a rule, require the use of time-consuming and expensive procedures. All the chemical methods are referred to destructive methods of control and are not direct methods, because they include a chemical preparation of samples for analysis. They are not also regarded as “multicomponent” methods allowing one to determine at once a practically complete collection of chemical elements contained in an object under study, as is possible in the case of the application of a number of spectral methods of analysis. Nevertheless, despite the difficulties mentioned above, the role of chemical methods remains significant, as usual, in the solution of special analytical problems. For example, in the case of an inhomogeneous distribution of trace amounts of a chemical element, its chemical separation and concentration allow one to eliminate the hampering influence of other components, to determine its lesser concentrations, and to overcome the difficulties related to both the inhomogeneous distribution of the required element in a sample and the absence of reference samples (standards). Such a problem has been solved by a chemical method in the present study.

In Sec. 8.2.1, we showed with the use of local methods of analysis (see Table 8.15) that the target explosion products contain, besides the common chemical elements, other chemical elements sufficiently rarely occurring in both, the Nature and the used structural materials. Among these rare chemical elements, we mention, for example, such chemical elements as beryllium, scandium, gallium, arsenic, selenium, strontium, yttrium, indium, iodine,

cesium, barium, thorium, uranium, and elements of the group of lanthanides. We are interested in, of course, not the very fact of the presence of the mentioned chemical elements in the explosion products. In fact, one can find any stable chemical element from the Periodical table in any specimen, though it will be present in an negligibly small amount. To prove that nuclear reactions are occurring upon the explosion of a target material, we must correctly demonstrate that the indicated process is accompanied by the increment of the amounts of the chemical elements, such as admixtures in the initial materials of articles participating in the explosion. We note that the establishment of this fact does not require us to solve the indicated problem for all chemical elements. Moreover, the solutions of this problem for different chemical elements have quite different convincing forces.

For example, if we solve the problem of widespread chemical elements, some doubts will always remain regarding the mentioned increment of their amounts resulting from the infertility upon the execution of experiments or the overlooking of certain inevitable sources of contaminations. At the same time, these doubts disappear to a significant degree, if a study is carried out as for chemical elements with negligible content in the initial materials of structural articles participating in the explosion. In the last case, they can appear in the explosion products only as a result of the nucleosynthesis. To such chemical elements, we would prefer to refer, first of all, to lanthanides. Their content in structural materials in those cases where they were not introduced specially as alloying admixtures (aluminum alloys, high-strength cast irons, and alloys for a special use) is of an extremely low level (see Refs. 133–135). As for stainless steel and sheet copper, it is well known that they do not contain lanthanides in the alloying composition.

It is obvious that to show the presence of chemical elements of the group of lanthanum in target explosion products with the use of local methods is a very simple task. However, to rigorously prove that their content increases upon the explosion of a target and to estimate the increment with satisfactory accuracy presents us with a rather complicated problem. This is related, first of all, to the circumstance that the distribution of these chemical elements in target explosion products is extremely inhomogeneous and their content is very low. To overcome the mentioned difficulties, we used chemical methods of analysis.

The study of target explosion products in order to detect the presence of lanthanides in them was carried out at A. V. Bogatsky Odesa Physico-Chemical Institute of the NASU. This group is well known by their investigations into elements of the group of lanthanum. As the initial specimens for analysis, we took copper accumulating screens with deposited products of the explosions of copper targets manufactured from the same material as the

screens (Cu of 99.99 mass. % in purity). Accumulating screens were of two types. We studied 17 screens of the first type. They were disks of 25 mm in diameter and 0.1 mm in thickness. Target explosion products were deposited on them only from one side. Screens of the second type were disks of 15 mm in diameter and 0.2 mm in thickness. Their number was equal to 25, and target explosion products were deposited on both their sides. The content of lanthanides in the specimens was determined by luminescence analysis with a luminescence spectrometer SDL-1 (see Ref. 137). Luminescence in the studied specimens was excited by means of the irradiation by ultraviolet from a gaseous-discharge mercury lamp known as a DRSh-250. Among all the lanthanides, the method used has revealed the highest sensitivity in the determination of the amounts of Eu and Tb. For their determination, we used the characteristic lines of the luminescence spectra of their ions with wavelengths of 612 nm [Eu(III)] and 545 nm [Tb(III)] conditioned by the transitions $^5D_0 \leftarrow ^7F_2$ and $^5D_4 \leftarrow ^7F_5$, respectively.

The procedure of qualitative analysis consists in the irradiation of the surface of accumulating screens with deposited target explosion products as received by ultraviolet and the registration of the characteristic luminescence of lanthanides. Upon the execution of the mentioned procedure, we registered only the weak luminescence of ions Tb(III). The qualitative analysis of the surface of accumulating screens without explosion products showed the full absence of any characteristic luminescence corresponding to chemical elements of the group of lanthanum. Thus, the results of the performed analysis showed that the layers of target explosion products deposited on accumulating screens contain only Tb in the registered amounts out of all the lanthanides.

To perform the quantitative analysis, we chemically separated lanthanides from the explosion products and concentrated them in a sample up to the enrichment level admissible for their registration (see Refs. 138–141). The indicated procedure consisted in that the accumulating screens with explosion products deposited on them were first dipped in a solution of hydrochloric acid (1 : 1) for 20 min. In this case, all screens of the first type were processed in the other solution than the screens of the second type, which led to both the washing-off of the surface layer from each accumulating screen with a mass of up to 0.2% to 0.3% of the screen mass and the accumulation of it in one of the solutions. Then, in both the solutions derived in such a way, we separated Cu from ions of the other metals. For this purpose, we introduced yttrium chloride (by 10 mg) in each solution and then added an aqua-ammonia solution. The deposit of formed yttrium hydroxide $Y(OH)_3$ (a nonselective collector), with which the ions of lanthanides were coprecipitated, was separated from soluble bright blue Cu ammoniates by filtering

with the washing on a filter. During the next stage, we separated the ions of lanthanides from the ions of other metals which coprecipitated with them on the nonselective collector. For this purpose, we dissolved the deposit of $Y(OH)_3$ in diluted hydrochloric acid. The subsequent addition of oxalic acid in the derived solutions led to the precipitation of the deposit of yttrium oxalate, which is a selective collector for lanthanides. Then yttrium oxalate with the oxalates of lanthanides captured by it was calcined at a temperature of 700°C in order to derive the oxides of lanthanides soluble in acids. To prepare the solutions for analysis, the oxides of lanthanides formed after the calcination were dissolved in hydrochloric acid. In the derived solutions, we determined the contents of lanthanides by the luminescence method (see Ref. 137). The quantitative analysis of Eu and Tb was realized with the use of their most sensitive analytical forms: the complex compounds with thenoyltrifluoroacetone and o-phenanthroline for Eu and the complex compound with nalidixic acid for Tb.

By using accumulating screens of the two mentioned types, we prepared two more solutions by an analogous scheme after the washing off their surface layers with the purpose of determining the content of lanthanides in the bulk of accumulating screens. On the first stage, we used diluted nitric acid (1 : 1) instead of hydrochloric acid for the quicker dissolution of the bulk of screens. All the subsequent operations, namely, the separation of Cu, coprecipitation of lanthanides on yttrium hydroxide and oxalate, preparation of the oxides of lanthanides by calcination, their dissolution in hydrochloric acid, and the luminescence-based determination of Eu and Tb, were performed exactly as in the previous case.

The results of determination of the content of lanthanides in the surface layers, being the deposited target explosion products, and in the bulk of accumulating screens are given in Table 8.22. This table contains the data concerning only Eu and Tb, because their presence was registered in all the solutions prepared for the investigation. As for the other lanthanides, Yb was registered in trace amounts only in two solutions derived from the washed surface layers of accumulating screens on the limit of sensitivity of the method. No other elements of the group of lanthanum were discovered in all the studied solutions. However, the last circumstance is possibly related to the fact that the employed method of analysis had a considerably lesser sensitivity regarding their (Sm, Dy, Tm, and Nd) determination (see Ref. 137).

Analyzing the presented data, we note, first of all, that both lanthanides are contained in considerable amounts in the surface layers of screens and in their bulk. In all the cases, their content exceeds the limits of detectability by the method used by several orders of magnitude, i.e., they

Table 8.22. Results of determination of the content of lanthanides in the surface layers and the bulk of accumulating screens.

Type of a screen	Total mass, g	Eu		Tb	
		$m_{\text{Eu}},$ μg	$C_{\text{Eu}},$ mass. %	$m_{\text{Tb}},$ μg	$C_{\text{Tb}},$ mass. %
Surface layers of screens					
1	0.0231	1.8	0.8×10^{-2}	6.0	2.6×10^{-2}
2	0.0122	7.3	6.0×10^{-2}	3.0	2.5×10^{-2}
Bulk of screens					
1	8.5674	2.8	3.3×10^{-5}	1.5	1.8×10^{-5}
2	5.6441	3.7	6.6×10^{-4}	12.5	2.2×10^{-4}

are reliably detected. We pay particular attention to the fact that the concentrations of Eu and Tb reached the level of several hundredths of a percent in the studied surface layers of accumulating screens (i.e., in target explosion products). As for the concentrations of these chemical elements in the bulk of accumulating screens remaining after the removal of the surface layers, they are lesser by two to three orders of magnitude. However, the amounts of these elements in the removed surface layers and in the remaining bulk of accumulating screens turn out to be commensurate in terms of total mass.

The last circumstance indicates, more likely, the partial removal of target explosion products upon the washing of the surface layers from accumulating screens rather than a high level of the content of lanthanides in the initial material of accumulating screens. Indeed, the accumulating screens of the first type lost only 0.269% of their mass upon the washing of the surface layers from them, whereas those of the second type lost 0.215% of their mass. This corresponds to the removal of a layer of about $0.13 \mu\text{m}$ or $0.22 \mu\text{m}$ from each side of screens of the first or second types, respectively. The partial washing of target explosion products is indicated by the following fact. On screens of the first type, the total content of lanthanides turns out more in the surface layers, than that in the bulk of screens ($7.8 \mu\text{g}$ against $4.3 \mu\text{g}$). But on screens of the second type, we observed the inverse relation ($10.3 \mu\text{g}$ against $16.2 \mu\text{g}$). Such a situation arises because only one layer of target explosion products was partially washed from each screen of the first type. As for screens of the second type, two layers of target explosion products were partially washed, since the target explosion products were deposited on both sides of these screens. Therefore, it turns out that the bulk of each screen of the first or second type contains the remnants of one or two layers with the target explosion products.

Table 8.23. Comparison of the content of lanthanides in accumulating screens with deposited target explosion products to their content in the initial material of accumulating screens.

Type of a screen	Total mass, g	Eu		Tb	
		$m_{\text{Eu}},$ μg	$C_{\text{Eu}},$ mass. %	$m_{\text{Tb}},$ μg	$C_{\text{Tb}},$ mass. %
Surface layers together with the bulks of screens					
1	8.5905	4.6	5.3×10^{-5}	7.5	8.7×10^{-5}
2	5.6563	11.0	1.9×10^{-4}	15.5	2.7×10^{-4}
Initial material of screens					
1			$< 1 \times 10^{-6}$		$\sim 1 \times 10^{-6}$
2			$< 1 \times 10^{-6}$		$\sim 1 \times 10^{-6}$

The above assumption was confirmed after the determination of the content of lanthanides in the initial material of accumulating screens and targets. The quantitative analysis of the blanks of accumulating screens and targets not participated in the explosion-involved experiments was carried out according to the above-described scheme. We established that the concentrations of Eu and Tb in them were at most 1×10^{-6} mass. %, i.e., by one to two orders of magnitude less than their values derived by us for the bulks of accumulating screens which were used in the explosion-involved experiments.

In Table 8.23, we give the data on the total content of each lanthanide for both types of screens with deposited target explosion products and their content in the initial material of accumulating screens. As can be seen, the total content of lanthanides in screens of the first type is 1.4×10^{-4} mass. %, and that in screens of the second type is 4.6×10^{-4} mass. %. In view of the fact that the total content of lanthanides in the initial material of screens (targets were manufactured from the same material) was at the level of about 1×10^{-6} mass. %, it is easy to conclude that the concentration of lanthanides on a copper accumulating screen increases due to the deposition of the explosion products of a copper target by more than two orders of magnitude. This assertion follows from the results of analyses of the composition of screens of both types with deposited explosion products and gives convincing evidence that the observed lanthanides were created as a result of the nucleosynthesis occurring upon the explosions of targets.

Now we compare the absolute values of the total amounts of lanthanides collected by accumulating screens of both types. They are $12.1 \mu\text{g}$ for screens of the first type and $26.5 \mu\text{g}$ for screens of the second type. The ratio of these amounts is equal approximately to 0.46. Let us try to clarify

the reasons for that the accumulating screens of the second type captured twice as much lanthanides. To this end, we mention that an accumulating screen of the first type was seen from the head part of a target at a solid angle of approximately 0.63π . This value can be easily calculated since the head part of the target (the takeoff point for the target substance upon its explosion) was located in all the experiments performed at a distance of 5 mm from the plane of an accumulating screen on the axis of symmetry passing through its center. Screens of the second type were seen from the head part of a target at a solid angle of 0.45π . The number of screens of the first type was 17; hence, the total solid angle covered by them upon the dispersion of the target substance was equal to 10.71π . The number of screens of the second type was 25, but the layer of explosion products was deposited on their both sides. Therefore, the total solid angle covered by them under the capture of the substance of targets was equal to 22.50π . The presented data show that the total solid angle covered by screens of the second type is considerably greater. It is logical to assume that this circumstance is the probable cause that the screens of the second type collected the greater amount of lanthanides. The ratio of the total solid angle covered by accumulating screens of the first type to that of the second type is 0.48, i.e., it almost exactly coincides with the ratio of the amounts of the captured lanthanides.

Let us discuss the derived equality. First of all, it is logical to assume that the source of lanthanides is positioned somewhere in the experimental chamber of the setup, because they appeared on accumulating screens as a result of the execution of experiments in it. It is also plausible to consider that the location of the mentioned source of lanthanides was invariable in all explosive experiments performed with screens of the first and the second types. Then it follows from the geometric reasons that the above-presented equality should hold for the point, at which the source of lanthanides is located, i.e., the amounts of lanthanides collected by screens of both types must be related each to the other as the values of the solid angles covered by them upon the dispersion of a substance from the point under consideration. But the symmetry of the problem yields that only a single point, whose location satisfies this equality, exists in the volume of the experimental chamber. The above-executed calculations indicate that the location of this point is in the head part of the target.

For logical completeness, let us consider the case, where the source of lanthanides is extended and the lanthanides are deposited on the working surfaces of accumulating screens condensing from the volume of the experimental chamber. In this case the ratio of the amounts of lanthanides collected by screens of both types must be equal to the ratio of total areas of

the working surfaces of accumulating screens. In the considered experiments the latter ratio is about 0.94. The lack of coincidence of the values of the considered ratios points to that of the assumption about the extended source of lanthanides is contrary to the fact. In other words, we have one more strong argument in favor of the assertion that the source of lanthanides (and of other chemical elements) is from the zone of collapse of the target substance located in its head part. Because the lanthanides were absent there prior to the explosion, we conclude that they are created as a result of nuclear reactions induced by the impact of the "coherent electron beam".

Consider one more aspect of the data presented in Table 8.23. It is reasonable to consider that, in the sequential identical explosive experiments, the same amounts of some chemical element are created on the average. However, then the sequential explosive experiments must conserve the ratios of the amounts of different created chemical elements. Indeed, the data given in Table 8.23 derived on the great sample of experiments yield that the ratio of the amounts of Tb and Eu collected by accumulating screens in the explosion-involved experiments is practically identical for screens of the first and second types and is equal to about 1.5. On the one hand, this fact indicates that the data of chemical analyses of lanthanides are quite exact and sufficiently reliable. On the other hand, it confirms the earlier established absence of the separation of chemical elements upon the dispersion of the target substance (see Sec. 8.2.1). We also note that, in one explosion-involved experiment, a screen of the first type captures $0.27\text{ }\mu\text{g}$ of Eu and $0.44\text{ }\mu\text{g}$ of Tb. The similar values for a screen of the second type are 0.22 and $0.31\text{ }\mu\text{g}$, respectively.

Finally, we consider the question that has no direct relation to the chemical analysis of lanthanides, but to the aspect of the appearance of lanthanides in target explosion products. In Fig. 8.25, we present the X-ray spectrum registered recently by X-ray electron probe microanalysis from one of the particles contained in the explosion products of a lead target which were deposited on copper accumulating screen No. 130. In this spectrum, it is noticeable not only that it contains about 1.5 mass. % of one of the most rare lanthanides Tm (its content in the earth's crust is at most $8\% \times 10^{-5}\%$ [see Refs. 78, 133, 138, 139]), but also the chemical environment, in which this lanthanide is registered in the given spectrum.

In this spectrum, thulium is a single representative of lanthanides, and this fact is very astonishing. The point is in that lanthanides in the Nature are observed only in groups (families or earths). If Tm would find itself in the explosion products as a natural admixture, it must be detected there as a representative of the yttrium earth, as a minimum, together with dysprosium, holmium, and erbium (see Refs. 78, 133). We could assume that

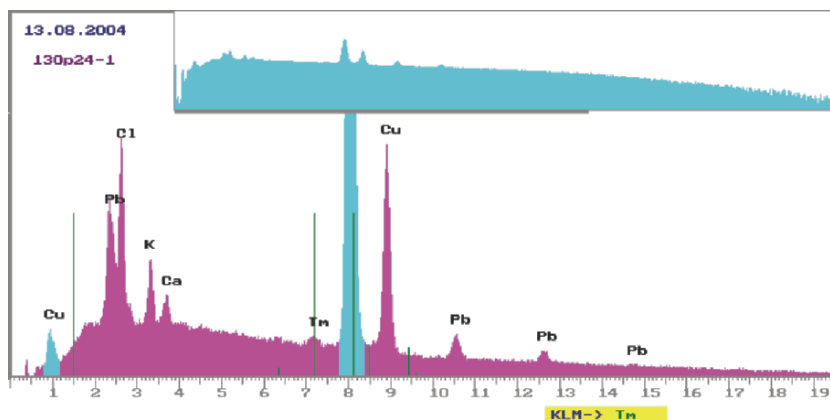


Fig. 8.25. Characteristic X-ray spectrum registered from a particle contained in the explosion products of a Pb target on Cu accumulating screen No. 130.

Tm got in the explosion products as an alloying admixture, but it should be noted that neither copper, not lead are alloyed with lanthanides (see Refs. 134, 135). Moreover, lanthanides as alloying admixtures are introduced in materials (aluminum and magnesium alloys, high-strength cast iron, high-quality optical glasses) usually in the form of a mischmetal consisting of a mixture of lanthanides of the cerium subgroup. If thulium is present in such a mixture, it is far from being its most predominate representative (see Ref. 134). In a more or less purified form, Tm can be introduced in microamounts in semiconductor materials, in particular, in gallium arsenide (see Ref. 133). But if it got into the explosion products together with any semiconductor, then it is not clear why there would be no gallium and arsenic in the spectrum. In other words, it seems to us that the most realistic explanation for the appearance of the detected thulium is its creation in the nuclear reactions which we believe occur upon the explosion of a target substance.

In conclusion, the administration of the Electrodynamics Laboratory “Proton-21” would like to thank the Director of A. V. Bogatsky Odesa Physico-Chemical Institute of the NASU, Academician S. A. Andronati and the staff of the Institute for their qualified help in the execution of the chemical analysis of lanthanides.

8.3. Main Results and Conclusions

We now summarize the studies of the isotope and element composition of target explosion products. At present, the main methods of determination of the isotope composition are various methods of mass-spectrometry. Here,

we used a number of methods and devices of mass-spectrometry, in particular: a mass-analyzer of gases, laser mass-spectrometry, glow-discharge mass-spectrometry, and secondary-ion mass-spectrometry. Data derived with the use of the methods of mass-spectrometry are not always simple for interpretation, which leads frequently to erroneous results and incorrect conclusions. The basic points in the derivation of correct results upon the application of these methods are the homogeneity of objects under study, the adequate decoding of complex spectra, and the identification of superpositions of the mass-peaks of complex and multiply charged ions with analytic mass-peaks.

In the present investigation, we have derived a huge amount of facts testifying to violations of the isotope composition of chemical elements in target explosion products under conditions when the above-mentioned difficulties were lacking. For example, anomalies of the isotope composition of inert gases were discovered on homogeneous objects (gas samples) upon the processing of simple mass-spectra containing no interferences. Upon the decoding of simple mass-spectra containing no peaks of clusters, we derived all the results by laser mass-spectrometry. Upon the determination of the isotope composition of basic chemical elements of a target material in products of the explosion by glow-discharge mass-spectrometry, the effects of superposition of mass-peaks were negligibly small. The last assertion is also completely true for results of the determination of the isotope composition of basic chemical elements of the materials of accumulating screens. As for the determination of the isotope composition of minor chemical elements of a target material by glow-discharge mass-spectrometry and secondary-ion mass-spectrometry, we took great care to use all the up-to-date correct procedures available in the scope of these methods to derive the above-presented results.

In other words, we understand that various results derived in this investigation are reliable to a variable degree, but the very fact of the establishment of anomalies of the isotope composition of chemical elements in target explosion products is indisputable for us, because the anomalies were registered many times in the solid and gas phases of these products and were confirmed by different methods. We note that this study has demonstrated the absence of both the fractionation of isotopes upon the preparation of gas samples and their separation upon the dispersion of a target material after the explosion. Both claims could be sources of the anomalies we observed, and needed to be investigated to eliminate them as the sources of the anomalies, which has been done. Therefore, we can conclude that the discovered anomalies of isotope compositions give a strong argument indicating the existence of significant numbers of nuclear reactions in the target

substances which occur upon its explosion induced by impact with “coherent electron bunches” from our experimental setup.

In the present work the element composition of target explosion products were studied with various highly sensitive physical (local and integral) and chemical methods of investigation. In all the cases it was established that, due to the shock compression of a target, the products of its explosion that are observed on an accumulating screen contain chemical elements that were not present earlier as impurities in the initial materials of both the target and accumulating screen in registered amounts. The last, as well as the discovered anomalies of isotope composition explosion products, is strong evidence of that we have observed artificially initiated intense processes of nucleosynthesis and transmutation of chemical elements in microvolumes of the substance of targets undergone to the explosion-induced compression up to superhigh densities, i.e., target explosion products contain nucleosynthesis products. Based on the results of these studies, we may conclude that the total mass of nucleosynthesis products originating upon the explosion of a target substance is on average 310 to 780 μg , which corresponds to 3.5×10^{18} to 8.0×10^{18} regenerated atoms or 1.7×10^{20} to 4.6×10^{20} nucleons.

Part III

Synthesis of Superheavy Elements in the Explosive Experiments

ON THE DETECTION OF SUPERHEAVY ELEMENTS IN TARGET EXPLOSION PRODUCTS

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Studying the composition of products of the laboratory nucleosynthesis by physical methods, in particular by glow-discharge mass-spectrometry, a number of experimental facts were established that indicate the presence of undetected chemical elements in addition to the known registered ones. The representatives of the mentioned undetected part of nucleosynthesis products are presumably stable, superheavy chemical elements positioned beyond the limits of the known part of the Periodic table of chemical elements. Moreover, the problem of search, registration, and study of representatives of this undetected part of nucleosynthesis products was posed. For its solution, we used methods of spectroscopy for studying the characteristic emission of different types generated upon a change in a state of electron shells of the atoms of chemical elements, methods of mass-spectrometry separating nuclides by their masses, and the purely nuclear method of Rutherford backscattering of α -particles and ions of N^{14} .

9.1. Discovery of X-Ray and Auger-Radiation Peaks from the Composition of Explosion Products

With the purpose of searching for superheavy chemical elements, we chose methods of the spectroscopy of characteristic emissions such as local Auger-electron spectroscopy and X-ray electron probe microanalysis. Such a choice was conditioned, first of all, by the fact that the mentioned methods have high locality and sensitivity in the determination of chemical elements (see Refs. 111, 120). It was obvious that these properties of the methods employed for the search for superheavy chemical elements are quite necessary. This followed from the previous investigations of target explosion products by the same local methods. These investigations revealed that the almost continuous layer of explosion products covering an accumulating screen consists mainly of the initial material of a target, and the nucleosynthesis products proper are embedded there in the form of small particles, drops, splashes,

and thin films. It became apparent that the search for superheavy chemical elements should be concentrated in such areas. We felt they would be either: distributed in the objects under study uniformly with low content or, on the contrary, be concentrated in small particles and thin submicron films. In both case, only the high locality and sensitivity of the methods of determination used for the composition could guarantee the detection of superheavy chemical elements. It should be noted that the above-mentioned methods of analysis have the range of determined chemical elements that is basically unbounded in respect to elements of large atomic numbers (see Refs. 111, 120).

9.1.1. Auger-Electron Spectroscopy

Local Auger-electron spectroscopy is a nondestructive method of quantitative determination of composition, and specimens studied by this method can be investigated by other ones. The method has a high spatial locality (50 to 100 nm), low depth of analyzed domain (1 to 2 nm), wide range of determined elements (all except for H and He), and high sensitivity (0.1 to 1 at. %). Upon using the ion etching, the method allows one to study also the distribution of a composition over depth (see Refs. 102–104, 119–122).

We searched for superheavy chemical elements in products of the laboratory nucleosynthesis on an Auger-microprobe JAMP-10S (the JEOL firm, Japan). Spectra were registered at an accelerating voltage of 10 kV on the electron probe, the beam current of 10^{-6} to 10^{-8} A, and the residual pressure of 5×10^{-7} Pa in the chamber with specimens. For the elimination of such artifacts as characteristic energy losses, we used accelerating voltages of 5 and 3 kV. The energy range of the semicylindrical mirror energy analyzer of an Auger-spectrometer was 30 to 3000 eV and the energy resolution was from 0.5% to 1.2%. All spectra were registered in differential form. For the quantitative analysis, we used the standard calculational program for the concentrations of elements supplied by the JEOL firm-producer of the device.

As objects for the investigation, we took exploded targets (wires of 0.5 to 1 mm in diameter, see Fig. 8.1a), that were manufactured from light, medium, and heavy chemically pure metals with atomic masses in the interval from 9 to 209, and accumulating screens with deposited products of the explosion of a target (see Fig. 8.1b). As their material, we usually took chemically pure V, Cu, Nb, Ag, Ta, and Au. Screens had the form of disks of 0.5 to 1.0 mm in thickness and 10 to 15 mm in diameter and served as a substrate. One of the surfaces of a screen was covered by the layer of a material that consisted of products of the nucleosynthesis. The layer revealed a weakly

pronounced relief, was characterized by axial symmetry, and its area varied in the range 1 to 2 cm². To determine the composition of a surface, we took specimens as-received. That is, the analyzed surface did not undergo any procedure of cleaning prior to its study that would destroy it or change its composition.

Below, we describe briefly the main strategy used in the search for superheavy chemical elements in nucleosynthesis products. On the surfaces of exploded targets and accumulating screen, we chose various micro-objects belonging to nucleosynthesis products, which can be judged by their morphology and, first of all, by their composition. From each micro-object, we registered the Auger spectrum with its subsequent interpretation. The essence of interpretation consists in trying to identify thoroughly all registered peaks by assigning them to the known chemical elements (see Refs. 123–126). Peaks which remained unidentified after such a decoding were verified for their belonging to artifacts of analysis (see Refs. 102–104, 119–122). The peaks we cannot associate with any known chemical element or refer to the artifacts of analysis were separated in a special group of basically unidentifiable Auger-peaks. Representatives of just this group of Auger-peaks were considered by us as candidates coming from superheavy chemical elements. In the described strategy, the basic links are the procedures of referring the observed Auger-peaks to the known chemical elements or the artifacts of analysis. By virtue of their importance, we dwell upon them in more detail.

The specificity of interpretation of the observed Auger-spectra consists in that they usually include Auger-peaks of more than 10 chemical elements. The collection of these elements can be considerably varied not only from specimen to specimen, but also from point to point in the limits of a specimen. Moreover, their contents were varied in a wide range. This led the spectra to include a great number of peaks which frequently overlap with one another. The mentioned circumstance and a number of other factors hampered the identification of Auger-peaks. Auger-spectra were usually registered in the energy range 30 to 3000 eV. The choice of a wide energy range for the registration of spectra was conditioned by two reasons. On the one hand, it allowed us to reach the processing of the maximum number of series of the Auger-peaks of analyzed elements and promoted the solution of the problem of their identification in the low-energy range. On the other hand, such a choice was necessary for solving the question about the presence of superheavy chemical elements in a specimen, because they, being “heavy” elements, have the most of their peaks in the high-energy range. However, the known reference books and catalogs of Auger-electron spectra (e.g., Refs. 123–126), usually used for the identification of Auger-peaks in

the registered spectra, contain an extremely bounded database in the energy range 1000 to 2000 eV. Moreover, they give no information about the energy range 2000 to 3000 eV.

Discussing the problem of identification of Auger-peaks, it is worth noting one more circumstance. The sensitivity of the method of Auger-electron spectroscopy in the determination of the amount of “heavy” elements is lower than that for “light” ones (see Refs. 119–122). Therefore, with the purpose of detecting small amounts of the required elements upon the registration of spectra and enhancing the signal-to-noise ratio, we used long-term exposures (up to three hours) and large currents of the primary beam of electrons (up to 10^{-6} A). However, the reference catalogs of Auger-electron spectra (see Refs. 123–126) used for the identification of the Auger-peaks of elements clearly give preference to high energy resolution (the record is realized on the narrow slits of a spectrometer), rather than to high signal-to-noise ratio. In other words, the reference books sacrifice low-intensity Auger-peaks to the fine structure of the analytic Auger-peaks of elements.

For the reasons clear from the above-presented discussion, we have found a number of undocumented Auger-peaks on specimens under study. These peaks are absent in the reference catalogs of Auger-electron spectra and, at the same time, belong to known chemical elements. The procedure of identification of each of these peaks was realized approximately by the following scheme. We analyzed the Auger spectrum, in which an unidentified peak was registered, and determined the collection of chemical elements, whose Auger-peaks were present in the analyzed spectrum. For each of these elements, we registered the standard Auger spectrum in the energy range 30 to 3000 eV with a high signal-to-noise ratio (long-term exposures) with the purpose of detecting undocumented low-intensity Auger-peaks. The standard Auger-spectra were registered on specimens from the corresponding pure elements or their simple compounds. Further, we tried to associate an unidentified peak with any low-intensity Auger-peak detected in the standard Auger-spectra of the chemical elements present in the analyzed spectrum. For example, we identified undocumented low-intensity Auger-peaks of the KLL series of Si and Al with an energy of 1737 and 1485 eV, respectively.

Some unidentified peaks were found to be the artifacts of analysis. As the artifacts of analysis, we considered such phenomena as the electric charging, characteristic energy losses, and chemical shift (see Refs. 99, 119–122). As a criterion for referring the unidentified peak to a peak of characteristic energy losses, we considered its shift or the disappearance upon a change in the accelerating voltage of the primary electron beam. About

electric charging, we judged by the correctness of the energy positions of the Auger-peaks of chemical elements present in the analyzed spectrum and its reproducibility upon changes in the accelerating voltage and the current of the primary electron beam. By analyzing such a phenomenon as chemical shift, we took into account, first of all, its value and the chemical environment of an analyzed peak. The peaks which then remained unidentified after the execution of the above-described procedures of analysis, their number being six, were referred to the class of basically unidentifiable ones. A number of these unidentifiable peaks has been accumulated as of the time of finishing the present work after the study of the composition of at least 100 specimens with nucleosynthesis products. Information about these peaks is given in Table 9.1.

The unidentifiable peak with an energy of 172 eV was registered on five specimens. It is contained in 11 Auger-spectra registered from these specimens. A fragment of one of such Auger-spectra containing the unidentifiable peak with an energy of 172 eV is shown in Fig. 9.1. The number of such spectra, if desired, can be arbitrarily increased, because a section of the surface of every specimen, where the spectrum can be registered, is sufficiently extended. The mentioned unidentifiable peak was registered on the surface of the metal matrix of an accumulating screen. Its intensity is low in all the cases and usually drops from the center of a screen to its periphery. The energy position of the peak is reproduced quite exactly in all measurements. Its typical chemical environment is Si, S, Cl, K, C, Ca, N, O, Cu, Zn, and Na, and nonmetal elements, except for C and O, are present usually in negligible amounts.

Some comments should be made on the stability of this peak. If, after the registration of the Auger-spectrum, we again register it at the same place, the intensity of this peak usually drops by 20% to 30%, whereas the subsequent repeated registrations show no noticeable decrease in its intensity. Moreover, we note that, in this case, no remarkable changes in the intensities of other peaks, in particular the Auger-peaks of C and O, occur, i.e., the indicated phenomenon cannot be referred to a deposit. It looks as if the substrate of this peak decays under the action of the electron beam of a probe in the near-surface layers, where it is most intense, and remains stable at great depths due to the weakening of the beam action. In the course of time, the intensity of this peak also decreases. This fact is shown in Fig. 9.2, where we present two Auger-spectra containing the unidentifiable peak with an energy of 172 eV that were registered at the same place, spectrum (b) being measured in after a period of one month. We note that spectrum (b) has the remarkably lower intensity of the unidentifiable peak than spectrum (a).

Table 9.1. Data on unidentifiable peaks in registered Auger-electron spectra.

Peak energy, eV	Typical element environment	Registration place	Behavior in time and under a probe	Specimen	N*
172	Si, K, Ca, Na, Cu, Zn, O, N, C, S, Cl	matrix	relatively stable	A132, A137, 4961, 5239, 5292	11
1095–1098	Si, K, Ca, Na, Cu, Zn, O, N, C, S, Cl	matrix	stable	A132, A137, 4961, 5239, 5292	11
525–528	Si, Al, Ca, O, C, S, Cl, Cu, Zn	particles	unstable, the content of Ca grows in time	4169, 4540	3
130 (115)	Al, O, C, N, S, P, Cl, Cu, Sn, Ce	matrix	unstable, the intensity drops in time	5633	2
94	C, O, Cu	particles, under the surface at a depth about 1 μm	stable	6215, a fragment of the shell	10
560	C, O, Cu	particles	unstable, the intensity drops in time	6754	1

*number of observations

On the description of the unidentifiable Auger-peak with an energy of 172 eV, it is necessary to indicate the circumstance that it appears always in the environment of the Auger-peaks of S and Cl with energies of 152 and 181 eV, respectively. Nevertheless, we do not associate the mentioned peak to such an artifact of analysis as chemical shift. On the one hand, S and Cl are present in the registered spectra always. On the other hand, even if this fact was interpreted as the chemical shift of the Auger-peak of Cl, it would

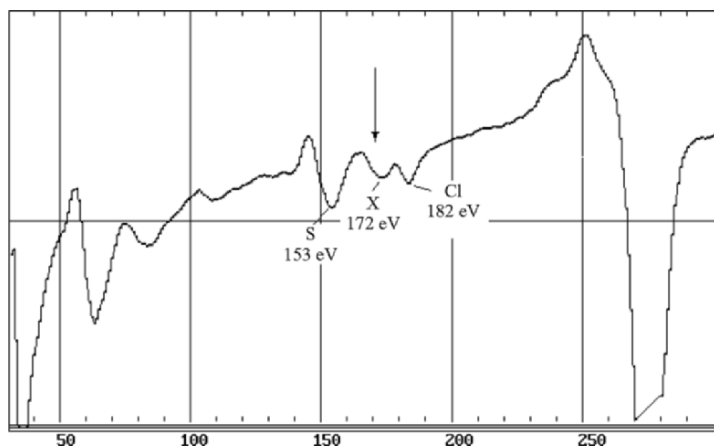


Fig. 9.1. Fragment of a typical Auger-spectrum containing the unidentifiable peak with an energy of 172 eV.

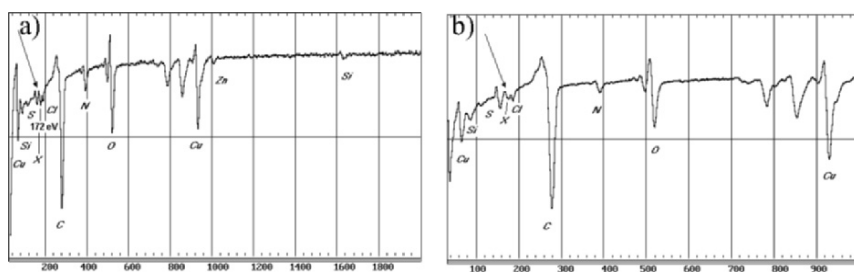


Fig. 9.2. Typical Auger spectrum containing the unidentifiable peak with an energy of 172 eV (a) and the Auger spectrum registered on the same place in one month (b).

be improbably large. In Table 9.2, we give the data on the most significant known chemical shifts of the Auger-peaks of S and Cl in their compounds.

Chronologically, the unidentifiable peak with an energy of 172 eV was discovered to be the first. Because we expected the presence of superheavy chemical elements in specimens and our main goal was their detection, it was reasonable to assume that the detected peak would be one of the most pronounced characteristic Auger-peaks of a certain superheavy chemical element, being present at the analyzed point of the specimen surface in a small amount. Then, in this case, the studied Auger-spectrum must also contain other Auger-peaks of some series of the same superheavy chemical element which would have a lower intensity. In other words, it was necessary to thoroughly study the whole energy range of the Auger-spectrum under consideration with the purpose of confirming the presence of low-intensity peaks in it.

Table 9.2. Data on the chemical shift of the Auger-peaks of S and Cl in some compounds (See Ref. 125).

Com- pound	$E(\text{Cl})_{\text{LVV}}$, eV	$\Delta E(\text{Cl})_{\text{LVV}}$, eV	Com- pound	$E(\text{S})_{\text{LVV}}$, eV	$\Delta E(\text{S})_{\text{LVV}}$, eV
LiCl	179	-2	PbS	149	-3
NaCl	182	1	Ag ₂ S	148	-4
CuCl ₂	178	-3	US _x	147.5	-4.5
KCl	178	-3			

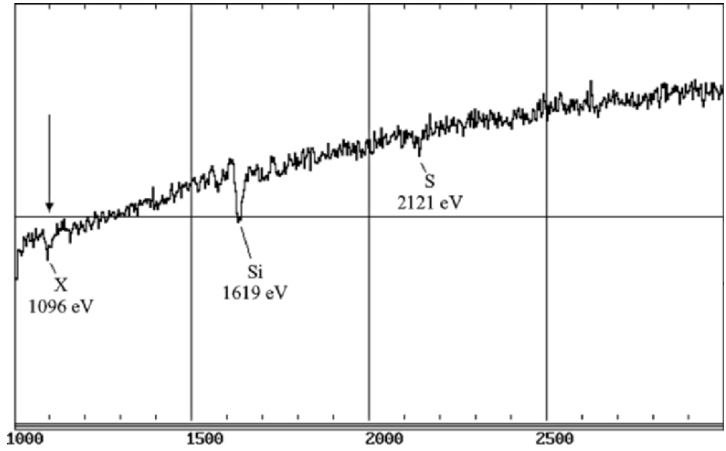


Fig. 9.3. High-energy part of the Auger-electron spectrum, whose low-energy part contains the unidentifiable peak with an energy of 172 eV. The exposure proceeded 3 h. The appeared unidentifiable peak with an energy of 1096 eV is regularly associated with the nonidentified peak with an energy of 172 eV.

The very first attempt of the registration of Auger-spectra using long-term exposure on those places of the specimen surface, where the unidentifiable peak with an energy of 172 eV was found, led to the discovery of a new unidentifiable peak with an energy of about 1096 eV. A high-energy fragment of one of these Auger-spectra is shown in Fig. 9.3. As for this new peak, it should be noted that it was registered without exceptions in all those cases where the unidentifiable peak with an energy of 172 eV was registered (see Table 9.1). That is, the former is regularly associated with the latter. By virtue of the fact that the intensity of the high-energy peak is always extremely small, it is rather difficult to mark its exact energy position due to the background of fluctuations. As a first approximation, we can consider it

to be in the interval from 1095 to 1098 eV. For the same reasons, the study of its behavior in time or under the action of the electron beam of a probe is a complicated problem. However, from the outside, it looks relatively stable.

The next unidentifiable peak had energy in the interval from 525 to 528 eV. A fragment of the Auger-electron spectrum containing this peak is shown in Fig. 9.4. This peak was registered in three Auger-electron spectra on two specimens (see Table 9.1). In all cases, the place of its registration was the surface of globular particles of the second phase (inclusions) with diameters up to 50 to 70 μm . Its intensity varied from low to significant (see Fig. 9.4). The peak appeared in the typical chemical environment: Si, S, Cl, C, Ca, O, Cu, Zn, and Al. The data on the chemical shift of the Auger-peak of oxygen are given in Table 9.3.

Considering the question on the stability of this peak, we will indicate several points. Its behavior under the action of the electron beam of a probe is such that, on the repeated registration of the Auger-spectrum following at once after the first one, it was not detected even in those cases where its intensity was considerable upon the first registration. Unfortunately, the Auger-spectra illustrating this fact were studied by an operator in the mode of fast scanning and were not registered, because they presented no interest at that time. However, two months later, we again returned to that particle, on which the intensity of the unidentifiable peak with an energy of 527 eV was maximum, and studied the composition of its surface. Both the corresponding Auger-spectra are shown in Fig. 9.5. They illustrate the fact

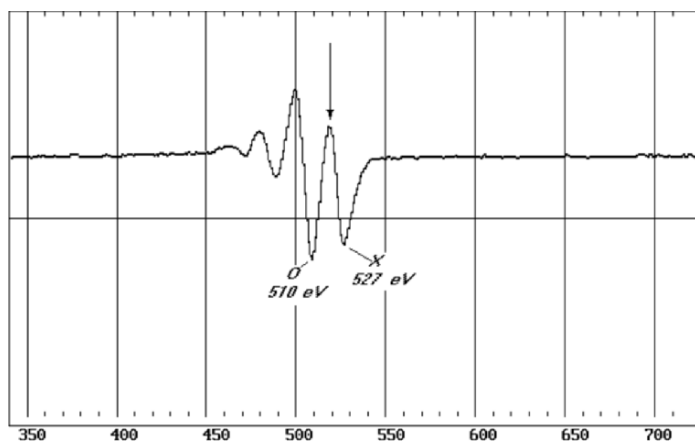


Fig. 9.4. Fragment of the Auger-electron spectrum containing the unidentifiable peak with an energy of 527 eV.

Table 9.3. Data on the chemical shift of the Auger-peak of oxygen in some compounds (See Ref. 125).

Compound	E(O) KLL, eV	$\Delta E(O)$ KLL, eV	Compound	E(O) KLL, eV	$\Delta E(O)$ KLL, eV
Mg(OH) ₂	503	-7	BeO	510	0
SrTiO ₃	503	-7	MgO	503	-7
α -Al ₂ O ₃	508	-2	Al(OH) ₃	511	1
SiO ₂	507	-3	Ca(OH) ₂	511	1
Ca(OH) _x	512	2	CaO	509	-1
MnO ₂	515.1	5.1	FeO	510	0
Fe ₂ O ₃	508	-2	NiOOH	511	1
Ni ₂ O ₃	512	2	CuO	502	-8
Cu ₂ O	509	-1	Y ₂ O ₃	507	-3
InPO _x	508	-2	Sb ₂ O ₅	507	-3
Nd ₂ O ₃	511	1	Dy ₂ O ₃	510	0
Tb ₂ O ₃	510	0	Tm ₂ O ₃	505	-5
Lu ₂ O ₃	510	0	HfO _x	510.8	0.8
PbO	513	3	SiN _x O _y	508	-2
LiNbO ₃	512	2	BaTiO ₃	509	-1
γ -2CaOSiO ₂	506	-4	[†] 3CaOSiO ₂	504	-6
β -2CaOSiO ₂	510	0	[†] 3CaOSiO ₂	511	1
CaOSiO ₅ 0.75TiO	512	2	[†] CaOSiO ₅ 0.5TiO	509	-1
CaO2Al ₂ O ₃	506	-4	[†] Ca ₃ Al ₂ O ₆	508	-2
[†] 3CaOAl ₂ O ₃	513	3	^{††} 3CaOAl ₂ O ₃	516	6
12CaO7Al ₂ O ₃	513	3	2CaOFeO	509	-1
4CaOAl ₂ O ₃ Fe ₂ O ₃	513	3	KAl ₃ Si ₃ O ₁₂	502	-8

[†]after hydration.

^{††}prior to hydration.

that the content of Ca on the particle surface grew considerably in the course of time.

Like the previous case for the unidentifiable peak with an energy of 527 eV, we undertook an attempt to detect the corresponding low-intensity peaks. However, upon the registration of Auger-spectra in a wide energy

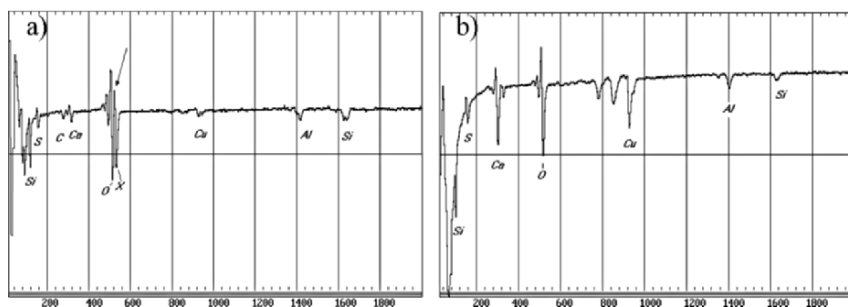


Fig. 9.5. Typical Auger spectrum containing the unidentifiable peak with an energy of 527 eV (a) and the Auger spectrum registered on the surface of the same particle in two months (b).

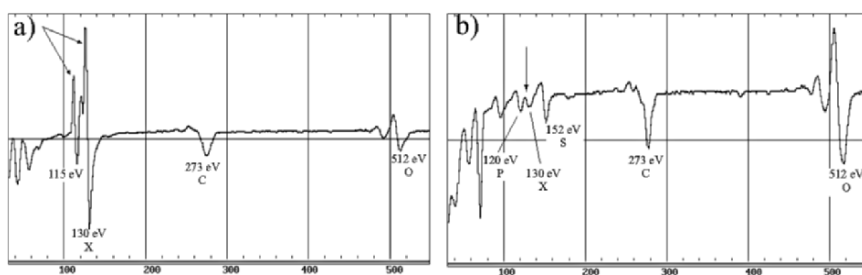


Fig. 9.6. Fragment of the Auger-spectrum containing the unidentifiable doublet of peaks with energies of 130 and 115 eV (a) and a fragment of the Auger-spectrum containing the unidentifiable doublet of peaks upon the repeated registration (b).

range and with a great time of exposure on the same places of specimens, where the unidentifiable peak with an energy of 527 eV was detected, we found no corresponding unidentifiable peaks.

The fourth unidentifiable peak (more exactly, a doublet of peaks) was registered only on one specimen. The energy of the major peak of the doublet was 130 eV, and that of the minor one was 115 eV. A fragment of the Auger-electron spectrum containing these peaks is shown in Fig. 9.6, a. The mentioned peaks were registered only in two Auger-electron spectra (see Table 9.1). In both cases, the place of their registration was small (5 to 10 μm in diameter) light-colored particles. The energy position of these peaks is reproduced quite exactly, and the intensity varies from a huge to a slight one. Their typical chemical environment is Al, O, C, N, S, P, Cl, Cu, Sn, and Ce.

Discussing the stability of these peaks, it is worth making some comments on their discovery. First, after the registration of a spectrum, they were erroneously identified as Auger-peaks of the MNN-series of yttrium,

because their energy coincides quite well. However, one circumstance was a puzzle: the registered peaks have high intensity, the spectrum was registered in the energy range from 30 to 3000 eV with long-term exposure, and the most intense Auger-peak of the LMM-series of yttrium with an energy of 1746 eV was absent. In this situation, the natural question arose: Must the intensity of the Auger-peak of the LMM-series of yttrium with an energy of 1746 eV exceed the background of fluctuations under the used conditions of the spectrum registration? To solve this problem, we registered the standard spectrum of yttrium from its simple compound Y_2O_3 under the same conditions. As a result, we established that the intensities Auger-peaks yttrium with energies of 127 and 1746 eV relate to each other 4 : 1. This means that if, in the situation under consideration, the low-energy doublet of peaks would belong to yttrium, then its Auger-peak of the LMM-series with an energy of 1746 eV should considerably surpass the background of fluctuations, because the intensity of the registered doublet was huge (see Fig. 9.6, a).

Having understood the fact that the discovered doublet of peaks does not belong to yttrium, we undertook attempts to register it again and to study its behavior and the existence area. In several days, we returned to the particle, where the doublet was first registered, but were unable to find it there. Sometimes, by studying a lot of similar particles, we managed to detect the studied doublet in the mode of fast scanning. However, in all the cases, its intensity was considerably lower than that on the first particle. Moreover, the doublet disappeared under measurements with long-term exposure except for one case. A fragment of the Auger-spectrum, in which the doublet was registered for the second time, is presented in Fig. 9.6, b. To the characteristics of the unidentifiable doublet, we refer the fact that it has no corresponding low-intensity unidentifiable peaks in the energy range from 30 to 3000 eV. This conclusion follows from the fact that such peaks were absent in the spectrum, in which the high-intensity unidentifiable doublet was registered under long-term exposure.

The fifth unidentifiable peak with an energy of 94 eV was registered on two specimens in 10 Auger-spectra (see Table 9.1). One of such spectra containing the mentioned peak is shown in Fig. 9.7, a. In all registered spectra, its intensity was rather considerable, and the energy position was reproduced quite exactly. Its chemical environment was the same in all the cases and rather poor, namely C, O, and Cu. It is characteristic that the environment did not contain even S, Cl, and N, that are usually present in all spectra on the specimens under study. In the course of time and under the electron beam of a probe, the peak was relatively stable.

Below, we will discuss the existence area of the unidentifiable peak with an energy of 94 eV. The specimen, where it was registered, was studied

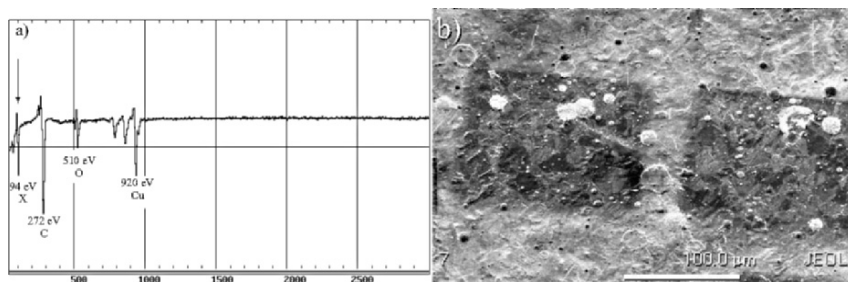


Fig. 9.7. Typical Auger-spectrum containing the unidentifiable peak with an energy of 94 eV (a); light-colored particles in craters from the ion etching, on which the peak was registered (b).

twice. From the very beginning, all the specimen surface and its details were investigated very thoroughly, but the peak was not detected. For the second time, the same specimen was studied approximately in a month during which it was studied by the method of secondary-ion mass spectrometry. We found a considerable number of craters of about $1\ \mu\text{m}$ in depth and $250 \times 250\ \mu\text{m}$ in area on the specimen surface that appeared after etching with the ion gun of a secondary-ion mass spectrometer. Both on the bottom of craters and on the specimen surface, we observed many round flattened light-colored particles of diameters from 10 to $40\ \mu\text{m}$ (see Fig. 9.7, b). These particles contained always the considerable amounts of Pb and, frequently, Si. We note that both elements contain rather intense low-energy Auger-peaks in their Auger-spectra in the range 92 to 94 eV. However, the spectra of Pb and Si contain, in addition, sufficiently intense high-energy Auger-peaks with energies of 2187 eV and 1619 eV, respectively, whose intensity is somewhat lesser. On some mentioned particles (only on those present on the crater bottom, rather than on the initial surface of the specimen), we registered spectra containing the peak with an energy of 94 eV and no Auger-peaks with energies of 2187 and 1619 eV. In this case, the intensity of this peak was so significant (see Fig. 9.7, a) that if it would belong to Pb or Si, their high-energy Auger-peaks would be revealed in spectra. By virtue of the fact that these Auger-peaks were absent in spectra, the 94-eV peak was referred to unidentifiable ones. The number of registrations of the unidentifiable peak with an energy of 94 eV can be arbitrarily increased, because it is found rather frequently. It is worth noting that the 94-eV peak has no associated unidentifiable low-intensity peaks in the energy range 30 to 3000 eV (see Fig. 9.7, a).

The next peak can be referred to unidentifiable ones only with certain reservations. It was registered only on one specimen and only in one spectrum (see Table 9.1). Its energy position can be estimated as 559 to 562 eV.

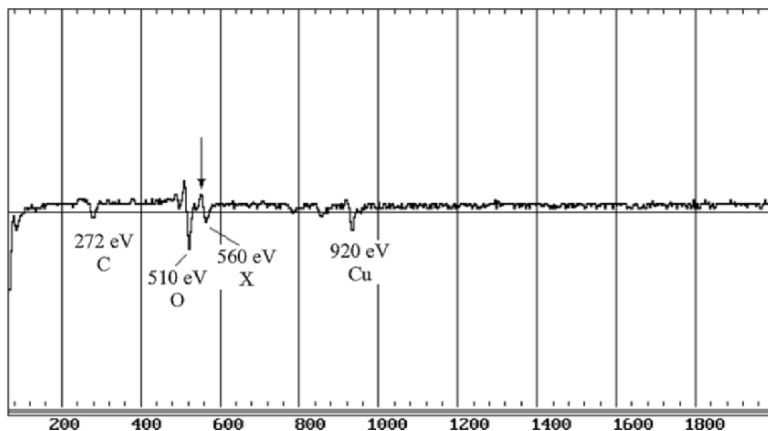


Fig. 9.8. Auger-spectrum containing the unidentifiable peak (Cs) with an energy of 560 eV.

The Auger-electron spectrum containing it is presented in Fig. 9.8. The place of its registration was small light-colored particle. Like the previous peak, it had a poor chemical environment, namely, C, O, and Cu. Its stability is characterized by the fact that the 560-eV peak was not detected at the repeated registration.

As for the reservations, we note the following. On the one hand, the energy position of the detected peak is such that it would be referred to the Auger-peak of the MNN-series of Cs with an energy of 563 eV, because the energies of the mentioned peaks are sufficiently close. On the other hand, it is obvious that the coincidence or proximity of energies of the peaks under consideration is not sufficient for such an identification. It is necessary also that the peaks of other series be present and coincide in energy. However, the other Auger-peaks of Cs have a slight intensity as compared to its analytical Auger-peak (563 eV), and the spectrum under study was registered under such conditions that the former would be not manifest above the background level (see Fig. 9.8). In other words, we have no sufficient reason to believe that this peak belongs to Cs, but, on the other hand, this assumption cannot be excluded. In any case, the behavior of the peak cannot be characterized as ordinary.

By summarizing the above-presented discussion of the derived experimental results, we note that, as one of the versions of the interpretation of the discovered unidentifiable Auger-peaks not referred to artifacts, it is natural, in the context of the present work, to assume their belonging to superheavy chemical elements. In this connection, of interest is the appraising determination of the atomic numbers of chemical elements which could be associated with some indicated Auger-peaks, e.g., the unidentifiable Auger-peaks with energies of 172, 527, and 1096 eV.

As one of the methods of estimation of the atomic numbers of the mentioned chemical elements, we can use the extrapolation based on the Moseley law (see Ref. 127). This law gives the dependence of the energy of a certain atomic level or the X-ray line energy for the fixed series of an atom on the atomic number. In view of our situation, the Moseley law can be written as

$$Z = k\sqrt{E} + Z_0, \quad (9.1)$$

where k and Z_0 are constants for one series of Auger-peaks, and Z is the atomic number of the chemical element, whose Auger-peak belonging to the series under study has energy E .

The data on the energies of the Auger-peaks of chemical elements can be found in reference books and catalogs (see Refs. 123–126). We are interested, first of all, in the energies of the Auger-peaks of chemical elements with large atomic numbers, but the corresponding information in the literature on this question is scanty. Analyzing the data given in Refs. 123–126, we can separate only two clear series of intense Auger-peaks of heavy chemical elements that are referred to Auger-transitions of the NOO type. The information on them is given in Table 9.4, where they are denoted conditionally as Series I and Series II. By Eq. 9.1 we determined the rated values of the atomic numbers Z' : $Z_{I0} = 58.00$ at $k_I = 2.48 \text{ eV}^{-1/2}$ and $Z_{II0} = 52.70$ at $k_{II} = 4.57 \text{ eV}^{-1/2}$ for Series I and Series II, respectively. The deviation of the rated value of the atomic number of a chemical element from its real atomic number, ΔZ , given in Table 9.4 testifies to that relation. Eq. 9.1 satisfactorily describes the series of Auger-peaks under study. We note that some increase in ΔZ with decrease in the atomic number is related to the growing error of measurements of the energy positions of Auger-peaks at small energies rather than to a deterioration of the degree of the used approximation.

Now, based on the derived constants (k_I , Z_{I0} , k_{II} , and Z_{II0}), we can estimate the assumed atomic numbers of the chemical elements which could be associated with the unidentified Auger-peaks with energies of 172, 527, and 1096 eV in the cases where the latter belong to Series I and Series II. The results of the mentioned calculations are given in Table 9.5.

We should like to make some comments on the derived data. First of all, it is seen from Table 9.5 that the Auger-peaks with energies of 172 and 527 eV could be referred to two different series of one chemical element with atomic number in the range 112 to 115. Moreover, we note as to the peak with an energy of 172 eV that its belonging to a chemical element with an atomic number of 90 to 91 seems improbable. Indeed, these elements have more intense Auger-peaks, but they are absent in the registered

Table 9.4. Approximation of two series of Auger-peaks from heavy chemical elements for an Auger-transition of the NOO type by the Moseley law.

Z	Element	Series I			Series II		
		Energy, eV	Z'	ΔZ	Energy, eV	Z'	ΔZ
92	U	188	92.00	0.00	74	92.01	0.01
90	Th	161	89.47	-0.53	64	89.26	-0.74
83	Bi	101	82.92	-0.08	44	83.01	0.01
82	Pb	94	82.04	0.04	40	81.60	-0.40
81	Tl	84	80.73	-0.27	36	80.12	-0.88
80	Hg	78	79.90	-0.10			
79	Au	72	79.04	0.04			
78	Pt	65	77.99	-0.01			
77	Ir	54	76.22	-0.78	25	75.55	-1.45
76	Os	45	74.64	-1.36			
75	Re	34	72.46	-2.54			
74	W	24	70.15	-3.85			

Table 9.5. Extrapolation of the atomic numbers of chemical elements for two series of Auger-peaks of the NOO-transition.

Series I		Series II	
Energy, eV	Z'	Energy, eV	Z'
172	90.52	172	112.63
527	114.98	527	157.71
1096	140.07	1096	203.92

spectra. As for Auger-peak with an energy of 1096 eV, the atomic number of a chemical element that could be associated with this peak should be considerable despite the lower accuracy of extrapolation in the case under study.

Discussing the results of extrapolation of the atomic numbers of chemical elements for registered unidentified Auger-peaks, we note that they could be referred also to other series of Auger-peaks, whose energies go beyond the limits of the available data (e.g., Refs. 123–126) on the energies

of the Auger-peaks of “heavy” chemical elements. It is obvious that, in this case, the greater values of the extrapolated atomic numbers of chemical elements would be derived. However, we have no basis for such an extrapolation at present.

9.1.2. Methods of X-Ray Spectrum Analysis

The first stage of solving the problem of detection of superheavy chemical elements in the products of the laboratory nucleosynthesis by the methods of X-ray spectrum analysis is, obviously, the registration of unidentifiable characteristic X-ray peaks in the spectra measured on them. Such peaks would testify to that the nucleosynthesis products contain chemical elements that are beyond the known part of the Periodic table of chemical elements, i.e., superheavy chemical elements. Prior to the representation of the results of investigations aimed at the search for and the registration of unidentifiable characteristic X-ray peaks in the nucleosynthesis products by the methods of X-ray spectrum analysis, we will characterize the situation, in which they were started.

By the indicated time, we have already studied the chemical composition of more than 800 exploded targets and accumulating screens with deposited products of the nucleosynthesis by the methods of X-ray electron probe microanalysis and X-ray fluorescence analysis. On these specimens, we registered about 13 000 spectra. The direct purpose of these studies was somewhat different and consisted in the analysis of the traditional chemical composition of specimens. Nevertheless, we gave a sufficient attention to the search for unidentifiable characteristic X-ray peaks in these spectra, because we have already registered all six unidentifiable Auger-peaks by Auger-electron spectroscopy. However, even the thorough analysis of these spectra revealed no unidentifiable X-ray peak in all cases. Moreover, we tried many times to find unidentifiable X-ray peaks by X-ray electron probe microanalysis on those micro-objects-representatives of nucleosynthesis products, on which we registered the mentioned unidentifiable Auger-peaks. All these attempts failed.

In other words, the problem of search for unidentifiable X-ray peaks cannot be solved by the progressive accumulation of the data on the X-ray spectra registered on nucleosynthesis products. Their absence cannot be explained by the fact that we failed to choose a suitable place for analysis in all cases of the registration of X-ray spectra. In fact, six unidentifiable Auger-peaks were registered on a sampling whose power was less at least by one order: we studied 108 specimens and registered 1745 Auger-spectra. In this case, we were guided in the choice of a place of analysis by Auger-electron spectroscopy, as well as by X-ray electron

probe microanalysis, by the same ordinary principles. All micro-objects-representatives of nucleosynthesis products were divided into classes by distinguishing their signs (morphology, contrast in the atomic number, etc.), and spectra were registered on several representatives of each class.

Based on the above reasoning, we may conclude that the unsuccessful results of searching for registered X-ray peaks upon the study of the composition of nucleosynthesis products should not be explained by the insufficient size of the sampling. On the contrary, it was more than sufficient. This yields that, upon solving the indicated problem, it is necessary to use the methods and procedures of the registration of X-ray spectra which are adapted to the problem, rather than the traditional ones. That is, we should rely on the realized strategy of searching, rather than on chance.

In order to construct such a strategy, we will analyze, first of all, the reasons for the failure of the attempts to detect unidentifiable X-ray peaks in the X-ray spectra registered upon the study of the composition of nucleosynthesis products. It is obvious that all these unsuccessful attempts are conditioned by a single reason, namely, by a low intensity of unidentifiable X-ray peaks. It is for this reason that we were unable to separate them from the background of fluctuations by using the traditional procedures of the registration of X-ray spectra.

This does not contradict the results obtained by Auger-electron spectroscopy. We recall that unidentifiable Auger-peaks had high intensities, and we had no problems in distinguishing them from the background. Indeed, in the case of a homogeneous distribution of superheavy chemical element in a specimen, we should manage to register unidentifiable X-ray peaks, because the sensitivity of X-ray electron probe microanalysis in the situation under consideration is higher than that of Auger-electron spectroscopy by almost one order. However, if a superheavy chemical element is distributed inhomogeneously and at least one characteristic size of its aggregates is in the far submicron range or the nanorange, we arrive at a different situation. In such a case, the volume of an aggregate of a superheavy chemical element can be equal to tens of percents of the volume of a domain analyzed by an Auger-microprobe and, hence, can be easily registered by local Auger-electron spectroscopy. But the volume of the same aggregate will be equal to only several tenths or hundredths of a percent of the volume of a domain analyzed by an X-ray microprobe, because the latter exceeds the volume of a domain analyzed by Auger-microprobe by almost three orders (see Refs. 111, 119). It is obvious that, in this case, the intensity of the X-ray peak of a superheavy chemical element will be at the detection threshold level.

Thus, the above analysis yields that superheavy chemical elements must be in the nucleosynthesis products in the form of aggregates or chemical

inhomogeneities on the submicron scale, and their characteristic X-ray peaks should be searched in X-ray spectra among low-intensity peaks. It becomes clear, first, which measures should be undertaken upon the registration of the X-ray spectra on nucleosynthesis products in order to maximally enhance the probability of the detection of unidentifiable X-ray peaks and, second, what realized strategy of their search must be chosen. First of all, one must guarantee a high locality upon the registration of spectra, because it favors the high sensitivity of the method. It is necessary also to make every effort for increasing the signal-to-noise ratio in the registered spectra. In the first turn, this can be realized at the expense of a significant increase in the duration of exposures upon the registration of spectra, since the level of a useful signal grows proportionally to time, and the background of fluctuations is proportional to the square root of time (see Ref. 127). The contrast of X-ray peaks can be also increased by decreasing the level of the very background, since this results in the damping of its fluctuations.

We have listed the main procedures that are used upon the registration of X-ray spectra and favor the improvement of the separation of low-intensity X-ray peaks from the background of fluctuations. We now turn to the very important question about where low-intensity unidentifiable X-ray peaks should be search for. The word “where” means the choice of the energy range of an X-ray spectrum with the maximum probability of the detection of a unidentifiable X-ray peak, rather than the choice of a suitable place of analysis on the specimen surface that contains a superheavy chemical element (in this case, all remains as before: the classification of micro-objects, their sorting, and intuition). To answer the posed question, we consider Fig. 9.9 that presents, as a diagram, the energy ranges of X-ray peaks of the K-, L-, and M-series of chemical elements of the known part of the Periodic table (see Ref. 130).

First, we turn to the energy range of characteristic X-ray peaks of the K-series. It follows from the diagram that the X-ray peaks of superheavy chemical elements must have energies above 120 keV. We cannot register them in principle, because the serial X-ray installation used by us gives no possibility to excite X-ray peaks with an energy above 40 to 50 keV. We now consider the energy ranges for X-ray peaks of the L- and M-series. As seen, characteristic X-ray peaks of these series of superheavy chemical elements may expand from 17 and 3.5 keV, respectively. These energy ranges are quite suitable for both the excitation of X-ray peaks and their detection on the available facilities. Thus, the above discussion yields that the unidentifiable X-ray peaks associated with superheavy chemical elements can be in the energy range beginning from 4 to 5 keV.

However, the last does not mean that they must be searched in this energy interval. This is connected with the circumstance that different sections of the indicated energy range have essentially different possibilities for the separation of a low-intensity unidentifiable X-ray peak from both the background of fluctuations and the superposition of the “tails” of X-ray peaks of other chemical elements. Indeed, let us divide the energy range under consideration into two subranges with energies from 4 to 17 keV and above 17 keV, respectively. As seen from the diagram (see Fig. 9.9), the first subrange includes X-ray peaks of the K-series of all chemical elements contained between Sc and Mo, and the L-series of all chemical elements contained between I and U. With regard to the presence of many components in the objects under study, this results in that the first subrange, as distinct from the second one, will be continuously filled by overlapping X-ray peaks of the known chemical elements contained in the specimens (see Fig. 8.15). It is obvious that to search for low-intensity peaks is hopeless in such an analytic situation. We also note that the background level and, hence, the value of its fluctuations in the first energy range is greater by almost one order of magnitude than that in the second one, which also does not favor the detection of low-intensity X-ray peaks contained in the former. Thus, the answer to the posed question is obvious: unidentifiable X-ray peaks should be searched in the second high-energy subrange, i.e., in the energy interval above 17 keV.

As the main method of study of the nucleosynthesis products with the purpose of detecting unidentifiable characteristic X-ray peaks in them, we chose X-ray electron probe microanalysis due to its high locality, sensitivity in the determination of chemical elements, and rapidity of analysis (see Ref. 111). The X-ray electron probe microanalysis was performed on an X-ray microanalyzer REMMA-102 (SELMi, the town of Sumy, Ukraine)

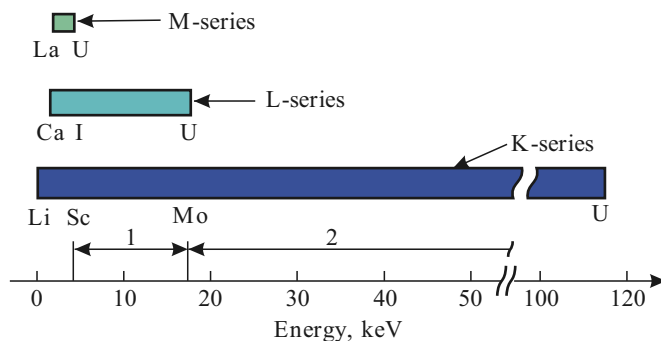


Fig. 9.9. Diagram of the energy ranges of X-ray peaks for the K-, L-, and M-series of chemical elements.

equipped with two-wavelength X-ray spectrometers and an energy-dispersive [Si(Li) detector] one. Spectra were registered at an accelerating voltage of the electron beam of 35 keV. The probe current and the electron beam diameter did not exceed 0.1 nA and 50 μ A, respectively. The residual pressure in the specimen chamber was not worse than 2×10^{-4} Pa. The energy-dispersive spectrometer ensured the registration of X-ray peaks in the energy range 0.3 to 32 keV, and its energy resolution on the line MnK_α (5.9 keV) was 160 eV. The used counting rate of X-ray quanta of 1200 to 1300 pulse/s guaranteed both the minimum “dead” time of the counting system and the practically full suppression of peaks of “sums” (see Ref. 111). The device was equipped with a COMPO detector allowing one to form the electron-scanning image of the surface of a studied specimen with contrast in atomic numbers of chemical elements contained in the specimen. The indicated mode was actively used for the choice of objects under analysis on the surface of the studied specimens.

With the purpose of expanding the possibilities of the registration of unidentifiable X-ray peaks in the high-energy range, we used, additionally, two setups for X-ray fluorescence analysis. One of them was equipped with an X-ray tube with Mo anode. The accelerating voltage on the tube was 45 keV. To excite secondary fluorescence, we used a section of the bremsstrahlung-spectrum of the X-ray tube with an energy of about 36 keV filtered by a monochromator made of pyrolytic graphite. The diameter of an analyzed domain on the surface of a specimen was usually at most 2 mm. The energy range of registered X-ray peaks was 1 to 35 keV.

In the second setup for X-ray fluorescence analysis, to excite secondary fluorescence, we used the soft γ -emission of an ^{241}Am isotope source emitting the intense line with an energy of 59.6 keV. The peak appearing in registered spectra due to the inelastic scattering of the source emission with maximum at an energy of 49 keV bounded the energy range of registered X-ray peaks from above by 43 keV. The diameter of an analyzed domain on the specimen surface in this setup was at most 1 mm. In both X-ray fluorescence setups, we used an energy-dispersive Si(Li)-detector for the registration of X-ray spectra. Like the case of X-ray microanalysis, its energy resolution on the line MnK_α (5.9 keV) was 160 eV. Upon the registration of X-ray spectra, the pulse counting rate on both setups for X-ray fluorescence analysis was at most several hundreds of pulses per second.

Below, we will discuss the used procedures of registration of X-ray spectra on the above-mentioned setups under searching for unidentifiable X-ray peaks. As was noted above, for their detection, it was necessary to use long-term exposures upon the registration of spectra. However, long-term exposures are usually accompanied by a drift of the working parameters

of devices: deviations of the accelerating voltage, displacements of the table with specimens, shifts of the probe from a point of analysis (in the case of X-ray microanalysis), variations in parameters of the amplifying channel of energy analyzers, etc. The mentioned phenomena can considerably affect the results of analysis. For example, a drift of parameters of an energy analyzer results in a shift of the energy scale, which, in turn, causes the broadening of X-ray peaks.

To prevent the undesirable effect of a drift of the device parameters on the results of measurements with the use of the method of X-ray fluorescence analysis, we employ the following procedure of registration of spectra. Spectra were registered in groups of 80 with the exposure duration for a separate spectrum of 1200 s. For each spectrum of a series, we carried out the calibration of the energy scale by means of its attachment to the reference points. As the last, we took the most intense X-ray peaks of the known energy contained in the calibrated spectrum. After the execution of such a procedure, all the spectra from a series were summed in a resulting spectrum with long-term exposure.

In the case of X-ray electron probe microanalysis, series consisted of 10 spectra with the duration of the exposure of a separate spectrum of 400 s. After the registration of each spectrum from a series, we carried out the control and correction of a position of the electron probe on the studied micro-object. Finally, calibrated spectra were summed into one resulting spectrum.

As the objects of investigation by the method of X-ray electron probe microanalysis, we used exploded targets and accumulating screens with explosion products of targets deposited on them. We described these objects in Sec. 8.2.1 and at the beginning of this section, where the results of registrations of unidentifiable Auger-peaks by the method of local Auger-electron spectroscopy are given. Therefore, we will not dwell on this theme in detail. We only recall that the electron-scanning micrographs of both a typical target after the explosion and a typical accumulating screen with target explosion products deposited on it are presented in Fig. 8.1, a, b.

For the method of X-ray fluorescence analysis, we used also specially manufactured specimens as the objects of investigation, besides the already-mentioned metallic accumulating screens with target explosion products deposited on them. The former were produced when a layer of target explosion products were washed away by diluted nitric acid (1:2) for 1 min from a typical metallic accumulating screen and deposited by evaporation on special light matrices-substrates transparent for X-ray emission. As such substrates, we used usually aluminum foils, aerosil powder, and various polymeric films. We used such specimens with a purpose of reaching a significant

suppression of the continuous background in the registered X-ray spectra at the expense of a decrease of the inelastic scattering of the primary emission by a light matrix. This was done to facilitate the separation of low-intensity unidentifiable X-ray peaks from the background of fluctuations.

For further presentation, it is necessary to refine the notion of unidentifiable characteristic X-ray peak. As an unidentifiable characteristic X-ray peak, we consider any peak contained in an X-ray spectrum registered on nucleosynthesis products that cannot be referred to the characteristic X-ray peaks of the known chemical elements (see Ref. 130) and is not the artifact of analysis. In other words, this is an X-ray peak, whose nature cannot be basically associated with any phenomenon or fact known at present, except for the characteristic X-ray emission of the chemical elements located beyond the limits of the known part of the Periodic table, i.e., superheavy chemical elements. As the artifacts of analysis, we considered the peaks of “sums”, peaks of losses, the background of fluctuations, the undocumented low-intensity characteristic peaks of the known chemical elements, and the peaks of elastically and inelastically scattered low-intensity undocumented lines of the soft γ -emission of the ^{241}Am isotope source (see Refs. 111–114).

We make some comments on the procedure of referring a registered X-ray peak to basically unidentifiable ones. In its first stage, an X-ray spectrum registered on nucleosynthesis products undergoes thorough analysis and interpretation. The procedure of interpretation of X-ray spectra consists in the identification of all characteristic X-ray peaks of the known chemical elements contained in it. The peaks remained to be unidentified, if it happens, underwent the thorough verification as for their belonging to the mentioned artifacts of analysis. And finally, the peaks passed through the filter of the verification on artifacts were referred to the class of basically unidentifiable characteristic X-ray peaks.

In the above-presented procedure, the most fundamental point is the stage of selection of else unidentified X-ray peaks by the criterion of their belonging to the artifacts of analysis. Therefore, we describe it in detail. We begin from such an artifact as peaks of “sums” (see Refs. 110–118). These false peaks appear in a spectrum when two X-ray photons come subsequently one after another into a detector during a shorter time interval than the characteristic time of its counting system. In this case, the counting system of the detector registers them as one X-ray photon with an energy corresponding to the sum of their energies. A peak of “sums” arises only if the described event occurs sufficiently frequently and with photons of the same energy rather than one time for the registration interval of a spectrum. The mentioned mechanism of formation of the peaks of “sums” yields that their appearance in a spectrum is favored by the

frequent entry of X-ray photons into a detector, i.e., by the high counting rates. It is also obvious that the most intense lines in the registered X-ray spectrum have the greatest probability to generate a peak of “sums”.

On the X-ray facilities used by us, the peaks of “sums” for the most intense lines become noticeable in spectra at the counting rates of about 2000 pulse/s. For this reason, we chose the counting rate on the X-ray microprobe by means of the tuning of the current of an electron beam so that its level be at most 1200 to 1300 pulse/s. On the installations for X-ray fluorescence analysis, the indicated condition was satisfied always because of their structural features.

The verification of an else unidentified peak for its affiliation to the peaks of “sums” was realized according to the following scheme. For all characteristic X-ray peaks with significant intensity contained in the spectrum under study, we constructed the table of the energy positions of the peaks of “sums” generated by them. If the energy of an unidentified peak is not equal to one of the energies from the mentioned table, the peak goes to the next stage of selection. If such a coincidence occurs, the fate of the unidentified peak depends on whether the peaks of “sums” generated by more intense spectral lines are revealed in the same energy range.

The other type of false peaks in X-ray spectra registered with energy-dispersive spectrometers is presented by the peaks of losses (see Refs. 110–118); sometimes they are called the takeoff peaks. Some words should be said about the mechanism of formation of these peaks. Upon the registration of an X-ray photon falling into a detector, its photoelectric absorption in Si occurs. The higher the energy of an incident photon, the more is the number of free charge carriers created by it in Si. By their number, the detector recognizes the energy of a registered photon. This energy will be correctly identified by the detector only in the case of the full absorption of a photon. The indicated process of absorption is multistage and rather complex. On separate stages, it is accompanied by the generation of X-ray K_{α} -emission of Si. If the SiK_{α} -photon appeared in such a way is not absorbed and leaves the detector due to a high penetrating ability of X-ray emission, the situation will be equivalent to the partial absorption of the initial photon. It is obvious that the detector underestimates, in this case, the energy of the registered photon by the energy of a SiK_{α} -photon (1.739 keV). If the described event occurs upon the registration of a spectrum with monoenergetic photons many times, the peak of losses will appear to the left from the peak corresponding to these photons at a distance of 1.739 KeV on the energy scale.

As usual, peaks of losses have negligible intensity, and one can discover them only near the most intense spectral lines. The test of an unidentified X-ray peak for its affiliation to peaks of losses is a very simple procedure.

It consists in the verification of the presence of an intense parent peak which must be located to the right from the verified peak at a distance of 1.739 KeV on the energy scale.

For undocumented X-ray peaks, we distinguish two types. The first type includes low-intensity characteristic X-ray peaks which belong to the known chemical elements and are not cataloged for some reason (apparently, due to their extremely low intensity). In order that these characteristic peaks be not erroneously considered as unidentifiable, we have detected the standard spectra of all chemical elements (Mo to Sm) such that the energies of X-ray peaks of their K-series fall into the studied energy interval 17 to 43 keV (see Fig. 9.9). All standard spectra were observed in specimens made of the indicated pure chemical elements or their simple compounds by the above-described procedure of registration using long-term exposures.

The other type of undocumented peaks in X-ray spectra includes the peaks of elastically and inelastically scattered nondescribed low-intensity lines of the soft γ -emission from an ^{241}Am isotope source (see Ref. 127). Due to their low intensity, these lines have not been studied, apparently at all, but sometimes they appear in observed spectra of the ^{241}Am . In order that they be not referred to as unidentifiable peaks, we must know their exact energy position. In essence, these peaks are the scattering peaks of the emission of a source and should appear in the observed spectra in pairs: each elastic peak should correspond to the inelastic peak located in the spectrum on the left from the former. Because the inelastic scattering cross-section of X-ray emission considerably exceeds the elastic scattering cross-section for heavy matrices, the elastic peak is usually revealed very weakly in this case (see Ref. 127). Therefore, to determine the exact energy position of these peaks, we used the scattering spectra of the emission of an ^{241}Am source for light matrices (organic glass). With the purpose of verifying the correctness of the identification of peaks, we calculated the energy shift of the inelastically scattered peak relative to the corresponding elastically scattered peak by the Compton formula (see Ref. 127) and then compared it with the observed shift in the registered spectrum.

The last from the considered artifacts is the peaks from fluctuations of the background (see Refs. 110–118, 127). The value of the background fluctuations can be arbitrarily large in theory. From the general viewpoint, the situation is possible where the most intense peak in the observed spectrum possesses the fluctuating nature. However, the probability of such a situation is negligibly small. For this reason, it is acceptable to consider peaks in the spectrum which have a considerable intensity as real peaks, rather than fluctuating ones, without any probabilistic estimations. In reality, it would be proper to consider every peak in the spectrum, including

peaks with considerable intensity, as a nonfluctuating one only with a certain probability, even if this probability is arbitrarily close to 1. The last means that low-intensity peaks are not worse, in principle, than peaks with considerable intensity. The former differ from the latter only by a somewhat lesser probability in that they are not fluctuating peaks.

In spectroscopy, the question about which probability is for a peak to be a nonfluctuating one, so that it can be referred as a real peak, is traditionally solved on the basis of the statistical “ 3σ ” criterion (see Ref. 127), where σ is the mean square background fluctuation. According to this criterion, a peak is considered to be real, if the probability for it to be a nonfluctuating one exceeds 0.9987 (we consider only positive fluctuations above the background level) or, respectively, if the probability for it to be a fluctuating one is less than 0.0013.

Thus, it follows from the above discussion that we must be able to estimate the probability for else unidentified peaks to be fluctuating peaks. Below, we demonstrate a procedure for the calculation of such estimates. The mean square background fluctuation σ is defined by the formula (see Refs. 142–151)

$$\sigma^2 = \frac{1}{N-1} \sum_{i=1}^N (I_i - \bar{I}_i)^2, \quad (9.2)$$

where N is the number of channels of the spectrum near the peak under study, in which the background level is measured; I_i is the measured intensity of the background in the i -th spectrum channel; $\bar{I}_i$ is the approximated mean intensity of the background in the i -th spectrum channel. To determine $\bar{I}_i$ on the spectrum section near the peak under study in 10 to 15 channels to the left and to the right from it, we approximated the mean level of the background with a power or linear function. The value of $\bar{I}_i$ corresponds to the value of this function in the i -th channel.

To determine the probability of a deviation of the background level in a separate channel by a given value, we consider the background spectrum as a collection of independent random values I_i obeying the normal distribution, where $i = 1, \dots, N$ is the channel number. Then the distribution probability density for any of them is described as (see Refs. 144–147)

$$\rho(I_i) = (1/\sqrt{2\pi}\sigma) \exp[-(I_i - \bar{I}_i)/2\sigma^2]. \quad (9.3)$$

The probability for the background level in a given channel to fall into some interval of values can be found by integration of the probability density (Eq. 9.3) over this interval. If deviations of the background are measured in units of σ , then the probability of the deviation of the background in the i -th channel from the mean value to the positive side by $n\sigma$ is (see Refs. 148–151)

Table 9.6. Probabilities of deviations of the background in one channel by a value exceeding $n\sigma$

n	$p_i(n)$
0.5	0.309
1	0.159
1.5	0.067
2	0.023
2.5	0.006
3	1.3×10^{-3}
3.5	2.3×10^{-4}
4	3.1×10^{-5}
4.5	3.4×10^{-6}
5	2.9×10^{-7}

$$p_i(n) = \int_{\bar{I}_i + n\sigma}^{\infty} \rho(I_i) dI_i = \frac{1}{2} \left[1 - \text{Erf}(n/\sqrt{2}) \right], \quad (9.4)$$

where $\text{Erf}(z)$ is the so-called error function integral, whose values are tabulated. In Table 9.6, we give the calculated values of the function $p_i(n)$ for some n .

Since a real peak cannot be represented by one channel in the spectrum (even low-intensity peaks contain at least 5 channels), we must be able to estimate the probability of the appearance of a fluctuating peak of the prescribed configuration, including its width and form. The configuration of any peak can be set by an ordered collection of numbers

$$\{n_1, \dots, n_k\}, \quad (9.5)$$

where k characterizes its width measured by the number of channels, and the ordered sequence of numbers n_i defines its form. Since the background spectrum can be considered as a collection of independent random values I_i , the probability for the background fluctuations to exceed values $n_1\sigma, \dots, n_k\sigma$, respectively, in k channels of the spectrum is

$$p(n_1, \dots, n_k) = \prod_{i=1}^k p_i(n_i) \quad (9.6)$$

with regard to Eq. 9.3.

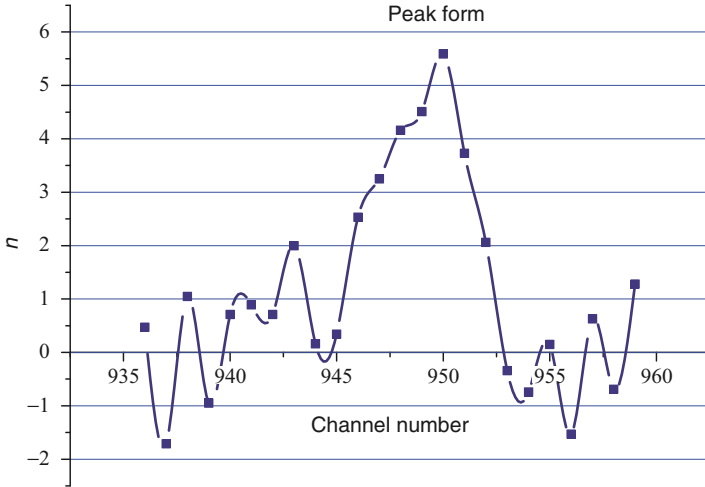


Fig. 9.10. Representation of an actual low-intensity peak as the intensity in units of the root-mean-square deviation σ versus the channel number for the X-ray spectrum.

Let us take, for example, an actual low-intensity peak observed in one of the spectra. Its configuration is presented in Fig. 9.10: the abscissa axis and the axis of ordinates show, respectively, the number of channels of the X-ray spectrum and the intensity in every channel in units of the root-mean-square deviation σ . In the numerical form, the data characterizing the given peak are presented in Table 9.7. In its second column, we give the experimental values of n_i , being the ratios of the intensity in the i -th channel to the root-mean-square deviation σ of the background. The third column contains the nearest minimum values of $n_{i,\min}$, for which we calculated the probabilities of deviations of the background (see Table 9.6). According to Eq. 9.6, the probability of the appearance of background fluctuations corresponding to the analyzed configuration of a low-intensity peak in 7 arbitrary channels is 1.26×10^{-27} .

The above-derived probability of the appearance of fluctuating deviations was calculated for 7 channels without regard for the ordering of the values of deviations of the intensity in channels. It follows from the experiment that the distribution of intensities in the adjacent channels for a real peak should have a clearly pronounced maximum and should be approximately normal. In other words, the given configuration Eq. 9.5 of a peak is not only a definite collection of intensities in k channels, but also the fixed order of their sequence. Since a peak is described by a single configuration and the number of possible configurations from k channels is $k!$, a proper accounting of the ordering of channels in the given configuration requires

Table 9.7. Configuration of an actual low-intensity peak from seven channels of the spectrum and the probability of the appearance of background fluctuations corresponding to its configuration in seven arbitrary channels.

Channel number	$n_i = I_i/\sigma$	$n_{i,\min}$	$p_i(n_{i,\min})$
946	2.53	2.5	0.006
947	3.25	3	$1.30 \cdot 10^{-03}$
948	4.16	4	$3.10 \cdot 10^{-05}$
949	4.51	4.5	$3.40 \cdot 10^{-06}$
950	5.59	5.5	$2.90 \cdot 10^{-07}$
951	3.73	3.5	$2.30 \cdot 10^{-04}$
952	2.06	2	0.023
$p(n_{1,\min}, \dots, n_{7,\min})$			$1.26 \cdot 10^{-27}$

multiplying probability Eq. 9.6 by $1/k!$. Because a fluctuating peak can appear at any place of the spectrum, we must multiply probability Eq. 9.6 by the number of its possible dispositions in the spectrum, i.e., approximately by the number of channels in the spectrum which is equal to 1024. As a result, the probability of the appearance of a fluctuating peak for the considered configuration of seven channels at an arbitrary place of the spectrum is about 2.5×10^{-27} . This value is considerably less than 0.0013 set by the “ 3σ ” criterion. The last means that the probability of the appearance of a peak for the considered configuration at the expense of background fluctuations is negligibly small, i.e., this peak should be considered as real.

Due to the smallness of the derived values and for convenience, we will use the rougher formula for the probability of the appearance of a fluctuating peak for the given configuration:

$$P(n_1, \dots, n_k) = 1000 \times p(n_1, \dots, n_k). \quad (9.7)$$

As compared to the exact formula, we replaced the number of channels, 1024, by 1000 and drop the factor $1/k!$ which is, of course, always less than 1.

To get the idea of the smallness of the intensity of an observed peak such that it can be else considered as a nonfluctuating peak, we present an evaluation of the values of the probabilities of the appearance of fluctuating peaks for some simple configurations and the numbers of spectra, which should be registered in order that the fluctuating peaks for the indicated configurations appear at least one time, in Table 9.8. We note that, among

Table 9.8. Probability of the appearance of fluctuating peaks for some configurations and the number of spectra which should be registered for the appearance of fluctuating peaks for the indicated configurations.

Peak configuration	$P(n_1, \dots, n_k)$	$1/P(n_1, \dots, n_k)$
1,3,1	3.29×10^{-2}	30
1,1,3,1	5.23×10^{-3}	191
1,1,3,1,1	8.31×10^{-4}	1204
1,3,3,1	4.27×10^{-5}	23406
1,2,3,2,1	1.74×10^{-5}	57518
3,3,3	2.20×10^{-6}	455166
1,3,4,3,1	1.32×10^{-9}	7.55×10^8
1,3,5,3,1	1.24×10^{-11}	8.07×10^{10}

the considered configurations of low-intensity peaks, only two configurations $\{1,3,1\}$ and $\{1,1,3,1\}$ (see the first and second rows of Table 9.8) do not satisfy the “ 3σ ” criterion. Thus, the given data yield that all low-intensity peaks reaching the level of at least 3σ at the maximum and possessing the width of more than four channels are not fluctuating peaks.

Thus, we have described the procedure of selection of the observed and unidentified X-ray low-intensity peaks as for their explanation as artifacts of analysis. As noted above, those unidentified peaks, which passed through the filter of verification for artifacts, were referred by us as a class of basically unidentifiable characteristic X-ray peaks. The data on all unidentifiable peaks observed as of March 2004 are presented in the summary Table 9.9.

An unidentifiable X-ray peak with an energy of 20.0 keV was observed twice by X-ray electron probe microanalysis (XEPMA). First, it was discovered on a metallic particle of about $1\mu\text{m}$ in diameter located on the surface of specimen No. 3, being a deposit of the nucleosynthesis products on the substrate made of Plexiglas. The spectrum containing this peak and its high-energy fragment are given in Fig. 9.11, a, b, respectively. The width of this peak is 11 channels, and it reaches a level of 4.6σ at the maximum. Its surrounding chemical environment is represented by such chemical elements as Pb, Cu, Au, Fe, Ni, Al, Cl, K, Ag, and Ca. Moreover, the X-ray peaks of the K-series of Cu possess the maximum intensity in the spectrum. Upon the repeated observation of the spectrum at the same point of analysis in 30 min after the first observation, we noticed a change in the composition. In the

Table 9.9. Data on unidentifiable peaks in observed X-ray spectra.

Peak energy, keV	Method of analysis	Typical element environment	Specimen	Expected ordinal number of the element		N^*	Peak height [†]
				by L_α	by M_α		
20.0	XEPMA	Pb, Cu, Au, Fe, Ni, Al Cl, K, Ca, Ag	No. 3 ¹ , No. 65 ²	109–110	198–202	2	4.6; 4.3
21.3	XFA with X-ray tube	Cu, Pb, Ag, Cd, Sn, In	Polymeric screen (film)	112–113	204–208	3	2.8
22.6	XEPMA	Pb, Cu, Fe, Ti, Ca, K, Si, Zn	No. 43 ²	115–117	209–214	1	5.5
23.6	XEPMA	Al, Cu, Pb, Ag	No. 8545 ² , No. 41 ² , No. 57 ²	117–118	213–218	4	4.5; 3.7; 3.5; 3.2
26.7	XEPMA	Cu, Al	No. 8545 ²	124–125	226–231	1	3.9
28.9	XEPMA	Ag, Pb, Bi, Sn, Cu	No. 8545 ² , No. 8858 ² , No. 41 ² , No. 43 ² , No. 47 ²	128–129	234–239	5	5.6; 4.6; 3.8; 3.2; 3.0
36.6	XFA with ²⁴¹ Am	Ag, Ba, Nd La, Pb, In Sn, Sb, Y Zr, Cd, I	No. 2 ³	141–143	260–267	1	5.3
38.1	XFA with ²⁴¹ Am	Ag, Ba, Nd	No. 1 ³	144–145	265–272	1	5.1
51.95	XFA with ²⁴¹ Am	Pb, Cu, Au, Fe, Ni	Polymeric screen (film)	162–164	306–314	1	4.4

* number of observations

† in units of σ

¹ deposit on organic glass

² copper screen

³ deposit on aerosil

repeated spectrum, the intensities of the peaks of Ca and Pb significantly decreased, and that of the peak of Cu grew. As for the very unidentifiable peak, it completely disappeared. We note the evidence of a change of the chemical composition and the disappearance of an unidentifiable peak upon repeated observations are the typical phenomena for many unidentifiable X-ray peaks we observed.

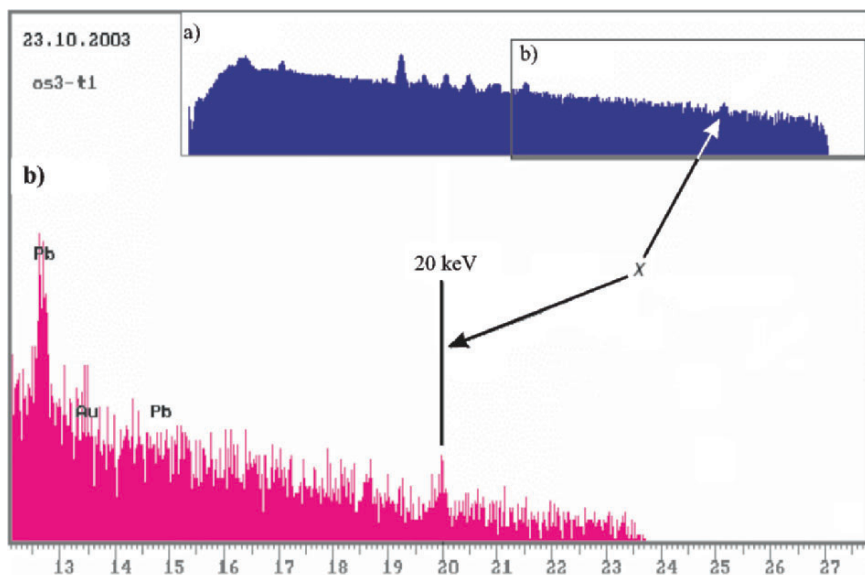


Fig. 9.11. X-ray spectrum containing the unidentifiable peak with an energy of 20.0 keV (a) and its fragment (b) presented on the logarithmic and linear scales, respectively. In part (a), fragment (b) is separated by a rectangular.

In this case, the disappearance of the unidentifiable peak can be caused by a movement of the analyzed particle due to the electric charging of the substrate made of Plexiglas or its partial melting because of the local heating by the current of the incident electron beam. However, in other cases where metallic screens were used as the substrate, similar reasons cannot explain the disappearance of unidentifiable peaks in the course of time. Moreover, their disappearance also cannot be explained by such known phenomena as the surface diffusion stimulated by an electron probe or screening by hydrocarbon films formed at the place of analysis due to the “carbon deposit”, because the temporal behavior of low-intensity peaks of the known chemical elements was ordinary and different from the unidentifiable peak behavior. Unidentifiable peaks behave themselves in the course of time as if their carrier (presumably a superheavy chemical element) undergoes the decay under the action of the electron beam.

The registered unidentifiable peak with an energy of 20.0 keV was not a peak of losses, because the intense parent peak to the right from it at a distance of 1.739 keV on the energy scale was absent (see Fig. 9.11). It cannot also be referred to the peaks of “sums”, because the analyzed spectrum did not contain the peak of “sums” with an energy of 16.09 keV

Table 9.10. Probability of the appearance of a fluctuating peak with the configuration of the unidentifiable peak with an energy of 20.0 keV.

Channel number	$n_i = I_i/\sigma$	$n_{i,\min}$	$p_i(n_{i,\min})$
641	0.479	0.5	$3.09 \cdot 10^{-01}$
642	2.239	2	$2.30 \cdot 10^{-02}$
643	3.209	3	$1.30 \cdot 10^{-03}$
644	0.936	0.5	$3.09 \cdot 10^{-01}$
645	4.355	4	$3.10 \cdot 10^{-05}$
646	2.516	2.5	$6.00 \cdot 10^{-03}$
647	1.202	1	$1.59 \cdot 10^{-01}$
648	1.200	1	$1.59 \cdot 10^{-01}$
649	2.948	2.5	$6.00 \cdot 10^{-03}$
650	0.927	0.5	$3.09 \cdot 10^{-01}$
$p(n_{1,\min}, \dots, n_{10,\min})$			$2.49 \cdot 10^{-17}$
$P(n_1, \dots, n_{10})$			$2.49 \cdot 10^{-14}$
$1/P(n_1, \dots, n_{10})$			$4.02 \cdot 10^{13}$

corresponding to the doubled energy of the peak of $\text{CuK}_{\alpha 1}$, being the most intense in the spectrum under consideration.

We present the probabilistic estimation that the unidentifiable peak with an energy of 20.0 keV is a fluctuating peak for the case of its second observation (Table 9.10) on specimen No. 65 which is a copper accumulating screen with deposited target explosion products. In the second case, its width was 10 channels, and the maximum intensity reached only 4.3σ (Fig. 9.12), being somewhat less than that in the first case. Nevertheless, the presented data (Table 9.10) yield that the probability for this peak to be a fluctuating one is about 2.49×10^{-14} , which is considerably lower than the critical level of the probability corresponding to the 3σ criterion. Moreover, it is necessary to record at least 4.02×10^{13} spectra for a single observation of the fluctuating peak for such a configuration. We note also that the unidentifiable peak under consideration is well approximated with the normal distribution (Fig. 9.12). Thus, the registered peak with an energy of 20.0 keV is a reproducible X-ray peak of the nonfluctuating nature, and it can be rightfully referred to as an unidentifiable one. The extrapolation of the atomic number of the chemical element, for which the considered

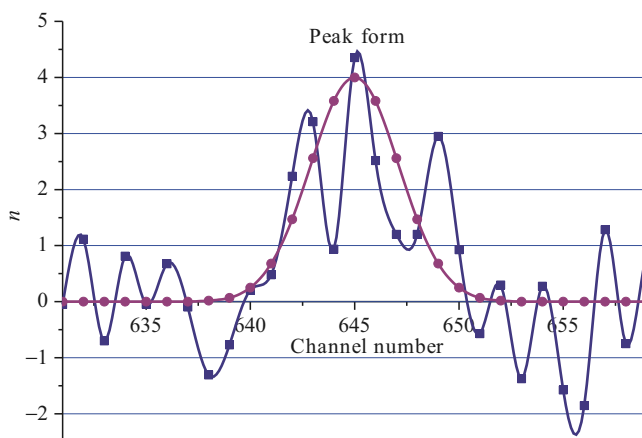


Fig. 9.12. Configuration of the unidentifiable peak with an energy of 20.0 keV as the dependence of the intensity in units of σ on the channel number of the X-ray spectrum (■) and its approximation with a normal distribution (●).

unidentifiable peak could be a characteristic peak belonging to the L- and M-series, by the Moseley law gives 109 to 110 and 198 to 202, respectively (see Table 9.9).

The X-ray peak with an energy of 21.3 keV was registered thrice on an installation for X-ray fluorescence analysis with an X-ray tube in the nucleosynthesis products deposited on an accumulating screen made of a polymeric film. The elements Cu, Pb, Ag, Cd, Sn, Sb, Zr, Nb, Mo, and In are its typical chemical environment. This peak was detected in the measurements at two points of analysis on the specimen. At one point, its appearance was confirmed by the repeated measurement executed in several days after the first discovery. However, upon the next observations of the spectrum at the same place of analysis, it disappears. One of the spectra containing the unidentifiable peak with an energy of 21.3 keV is given in Fig. 9.13. This peak was not an undocumented characteristic low-intensity peak of the known chemical elements. It was not referred also to the peaks of losses or “sums”. Nevertheless, we can refer it to unidentifiable peaks only conditionally since its maximum height in all the cases of registration was in the interval 2.5 to 2.8σ , i.e., it did not satisfy the “ 3σ ” criterion. At the same time, it is extremely improbable because of its reproducibility that it has the fluctuating peak nature. It could belong to the L- or M-series of chemical elements with numbers 112 to 113 and 204 to 208, respectively.

The unidentifiable peak with an energy of 22.6 keV was registered only one time by X-ray electron probe microanalysis on specimen No. 43

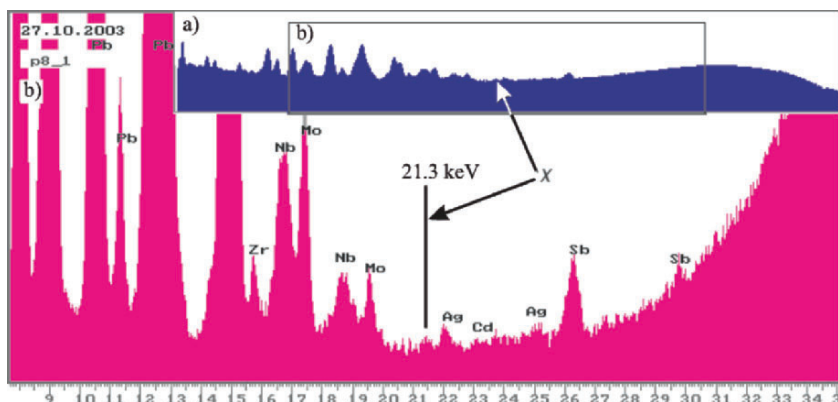


Fig. 9.13. X-ray spectrum containing the unidentifiable peak with an energy of 21.3 keV (a) and its fragment (b) presented on the logarithmic and linear scales, respectively.

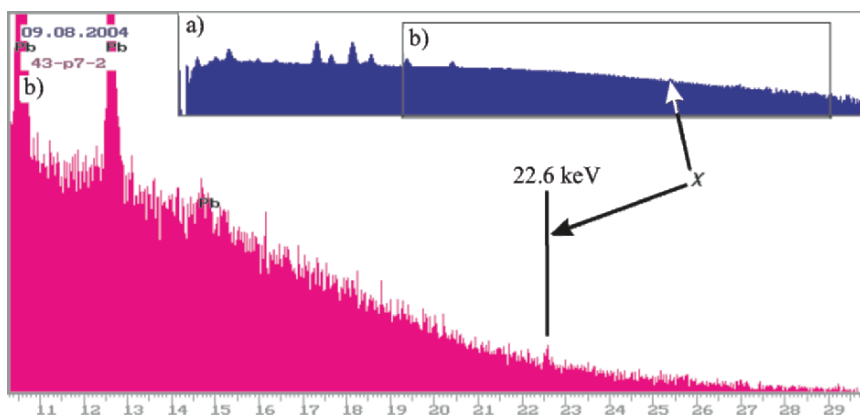


Fig. 9.14. X-ray spectrum containing the unidentifiable peak with an energy of 22.6 keV (a) and its fragment (b) presented on the logarithmic and linear scales, respectively.

(a copper accumulating screen). It was discovered in the spectrum recorded from a gray inclusion on the surface of a rounded particle of about $35\ \mu\text{m}$ in diameter. The spectrum containing this peak is given in Fig. 9.14. The peak was registered in the environment of such chemical elements as Pb, Cu, Fe, Ti, Ca, K, Si, and Zn. The X-ray peaks of the K-series of Cu and Fe and the L-series of Pb have the highest intensity in the spectrum. The registered peak had a rather high intensity: at the maximum, it reached the level of 5.5σ . The peak was well described by the normal distribution, and its width was five channels (see Fig. 9.15). It is not a peak of losses or “sums”. The peak

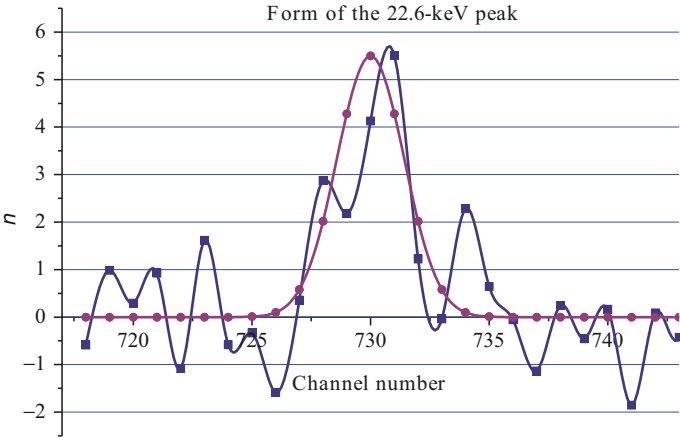


Fig. 9.15. Configuration of the unidentifiable peak with an energy of 22.6 keV as the dependence of the intensity in units of σ on the channel number of the X-ray spectrum (■) and its approximation with the normal distribution (●).

Table 9.11. Probability of the appearance of a fluctuating peak with the configuration of the unidentifiable peak with an energy of 22.6 keV.

Channel number	$n_i = I_i/\sigma$	$n_{i,\min}$	$p_i(n_{i,\min})$
728	2.874	2.5	$6.00 \cdot 10^{-03}$
729	2.179	2	$2.30 \cdot 10^{-02}$
730	4.128	4	$3.10 \cdot 10^{-05}$
731	5.509	5.5	$2.90 \cdot 10^{-07}$
732	1.224	1	$1.59 \cdot 10^{-01}$
$p(n_{1,\min}, \dots, n_{5,\min})$			$1.97 \cdot 10^{-16}$
$P(n_1, \dots, n_5)$			$1.97 \cdot 10^{-13}$
$1/P(n_1, \dots, n_5)$			$5.07 \cdot 10^{12}$

satisfies the 3σ criterion, because the probability for it to have a fluctuating nature was 1.97×10^{-13} (see Table 9.11). To register a fluctuating type peak with its configuration, we would record about 5.07×10^{12} spectra. According to the Moseley law, it could belong to the L- or M-series of chemical elements with numbers 115 to 117 and 209 to 214, respectively (see Table 9.9).

The unidentifiable peak with an energy of 23.6 keV was also registered by X-ray electron probe microanalysis. The number of its observations has

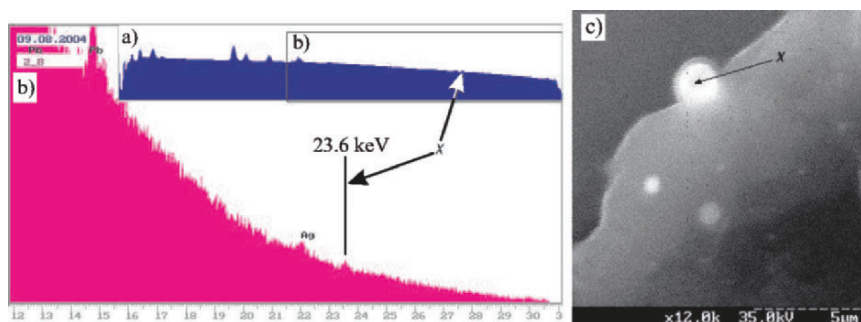


Fig. 9.16. X-ray spectrum containing the unidentifiable peak with an energy of 23.6 keV (a), its fragment (b) presented on the logarithmic and linear scales, respectively, and the scanning electron micrograph of a particle, on which the presented spectrum was registered (c).

reached four at present. It was twice discovered at different points of analysis on specimen No. 8545 and one time on specimens No. 41 and No. 57. All the mentioned specimens were Cu accumulating screens with deposited target explosion products. The maximum intensity of this peak varied from 3.2σ to 4.5σ . Like most unidentifiable peaks, it disappeared upon repeated observations of the spectrum from the place of its discovery. In Fig. 9.16, we present one of the spectra containing it and the scanning electron micrograph of a light-colored rounded particle of about $2\mu\text{m}$ in diameter, on which this peak was observed. The peak appeared usually in the chemical environment of Al, Si, Cu, Pb, and Ag.

Despite the closeness of the energy positions, this unidentifiable peak cannot be considered as a peak formed by the “sum” $\text{Cu } K_{\beta 1} + \text{Pb } L_{\alpha 1}$ with an energy of 23.665 keV, because the spectra, in which it was observed, did not contain the peaks formed by the “sum” and “doublings” corresponding to the peaks $\text{Cu } K_{\alpha 1}$ and $\text{Pb } L_{\alpha 1}$ (16.088 keV, 18.595 keV, and 21.102 keV), being more intense by one order of magnitude. The considered unidentifiable peak was not a peak of losses or an undocumented low-intensity characteristic peak of a known chemical element. In Fig. 9.17, we give the configuration of one of the observed unidentifiable peaks with an energy of 23.6 keV. It is symmetric, its width consists of 10 channels, and the intensity at the maximum is 4.5σ . The unidentifiable peak satisfies the 3σ criterion: the probability for it to have the fluctuating nature is 9.73×10^{-18} (see Table 9.12). It could be a characteristic peak of the L- or M-series of chemical elements with numbers 117 to 118 and 213 to 218, respectively (see Table 9.9).

The next registered unidentifiable peak has an energy of 26.7 keV (see Fig. 9.18). It was registered by X-ray electron probe microanalysis on

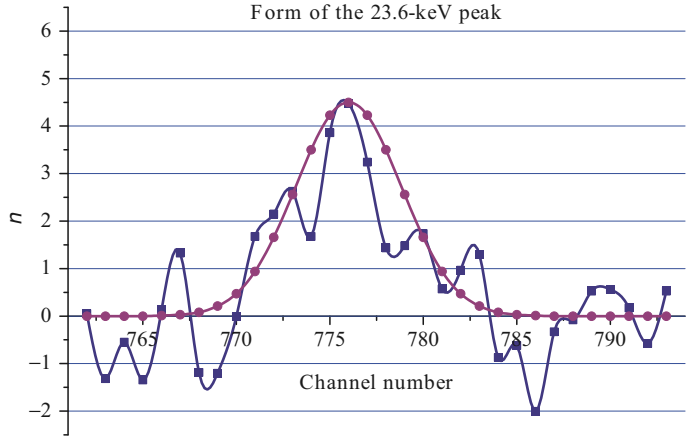


Fig. 9.17. Configuration of the unidentifiable peak with an energy of 23.6 keV as the dependence of the intensity in units of σ on the channel number of the X-ray spectrum (■) and its approximation with the normal distribution (●).

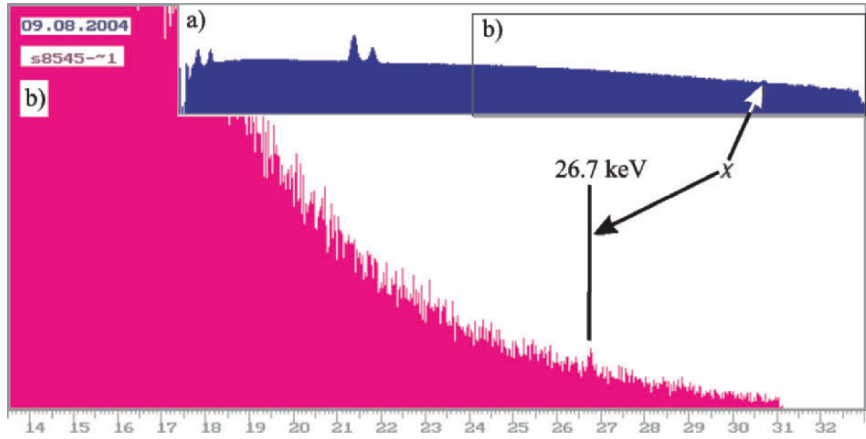


Fig. 9.18. X-ray spectrum containing the unidentifiable peak with an energy of 26.7 keV (a) and its fragment (b) presented on the logarithmic and linear scales, respectively.

the surface of a small dark spot on specimen No. 8545, i.e., on the same specimen, where the previous unidentifiable peak was recorded. We note that the peak was revealed in spectra recorded on this specimen (the screen and the target) else twice in other observations. However, its intensity in the two last cases did not exceed the level of 3σ . The peak had a rather poor chemical environment: Cu and Al. Its neighborhood did not contain

Table 9.12. Probability of the appearance of a fluctuating peak with the configuration of the unidentifiable peak with an energy of 23.6 keV.

Channel number	$n_i = I_i/\sigma$	$n_{i,\min}$	$p_i(n_{i,\min})$
771	1.675	1.5	$6.70 \cdot 10^{-02}$
772	2.152	2	$2.30 \cdot 10^{-02}$
773	2.627	2.5	$6.00 \cdot 10^{-03}$
774	1.679	1.5	$6.70 \cdot 10^{-02}$
775	3.861	3.5	$2.30 \cdot 10^{-04}$
776	4.477	4	$3.10 \cdot 10^{-05}$
777	3.242	3	$1.30 \cdot 10^{-03}$
778	1.436	1	$1.59 \cdot 10^{-01}$
779	1.481	1	$1.59 \cdot 10^{-01}$
780	1.737	1.5	$6.70 \cdot 10^{-02}$
$p(n_{1,\min}, \dots, n_{10,\min})$			$9.73 \cdot 10^{-21}$
$P(n_1, \dots, n_{10})$			$9.73 \cdot 10^{-18}$
$1/P(n_1, \dots, n_{10})$			$1.03 \cdot 10^{17}$

Table 9.13. Probability of the appearance of a fluctuating peak with the configuration of the unidentifiable peak with an energy of 26.7 keV.

Channel number	$n_i = I_i/\sigma$	$n_{i,\min}$	$p_i(n_{i,\min})$
869	1.93	1.00	$1.59 \cdot 10^{-01}$
870	2.01	2.00	$2.30 \cdot 10^{-02}$
871	3.91	3.50	$2.30 \cdot 10^{-04}$
872	3.00	2.50	$6.00 \cdot 10^{-03}$
$p(n_{1,\min}, \dots, n_{4,\min})$			$5.05 \cdot 10^{-09}$
$P(n_1, \dots, n_4)$			$5.05 \cdot 10^{-06}$
$1/P(n_1, \dots, n_4)$			$1.98 \cdot 10^{05}$

peaks of “sums”. It width was four channels, and it reached the level of 3.9σ at the maximum (see Fig. 9.19). The peak satisfies the 3σ criterion (see Table 9.13). It could belong to characteristic peaks of the L- or M-series of chemical elements with numbers 124 to 125 and 226 to 231, respectively (see Table 9.9).

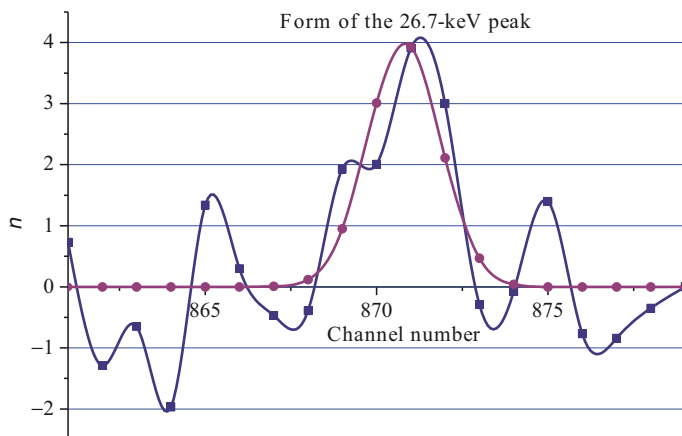


Fig. 9.19. Configuration of the unidentifiable peak with an energy of 26.7 keV as the dependence of the intensity in units of σ on the channel number of the X-ray spectrum (■) and its approximation with the normal distribution (●)

An unidentifiable peak with an energy of 28.9 keV is most frequently observed in the nucleosynthesis products. At present, it was detected in five cases. First, it was found by X-ray electron probe microanalysis on the same Cu screen No. 8545, where we detected the unidentifiable peaks with energies of 23.6 and 26.7 keV. It was registered on the surface of a rounded light-colored particle of about $15\ \mu\text{m}$ in diameter. In this case, it had a maximum intensity reaching the level of 5.6σ . The spectrum containing this peak is presented in Fig. 9.20, and its configuration is given in Fig. 9.10. The results of probabilistic estimations of the fluctuating peak nature are shown in Table 9.7 which yields that the peak satisfies the 3σ criterion. The peak was observed after on specimens No. 41, 43, 47, and 8858. Three first specimens were Cu accumulating screens, and the last one was a longitudinal microsection of the exploded target. On specimens No. 41 and 43, we observed earlier the unidentifiable peaks with energies of 23.6 and 22.6 keV, respectively. In all the cases of registration, the unidentifiable peak with an energy of 28.9 keV had a poor chemical environment (Ag, Pb, Bi, Sn, and Cu). For this reason, its nearest vicinity did not contain the possible peaks of “sums”. Upon the extrapolation by the Moseley law, the peak falls into the L-series and M-series of chemical elements with numbers 128 to 129 and 234 to 239, respectively (see Table 9.9).

The unidentifiable peak with an energy of 36.6 keV was observed on an installation for X-ray fluorescence analysis with the ^{241}Am isotope excitation source on specimen No. 2, being a layer of target explosion products

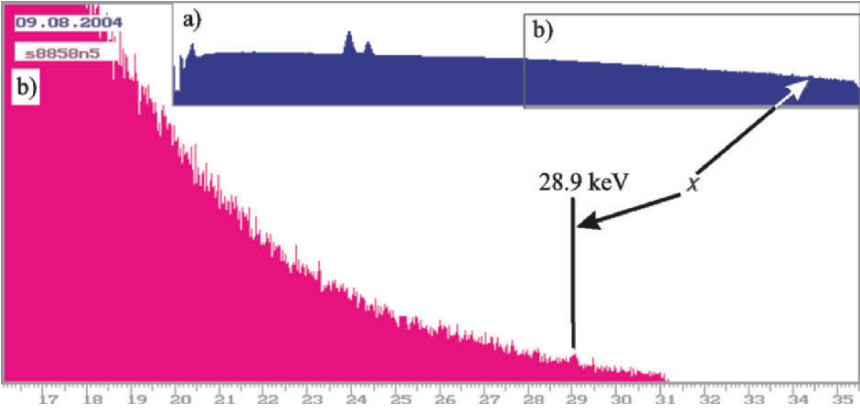


Fig. 9.20. X-ray spectrum containing the unidentifiable peak with an energy of 28.9 keV (a) and its fragment (b) presented on the logarithmic and linear scales, respectively.

Table 9.14. Probability of the appearance of a fluctuating peak with the configuration of the unidentifiable peak with an energy of 36.6 keV.

Channel number	$n_i = I_i/\sigma$	$n_{i,\min}$	$p_i(n_{i,\min})$
450	2.8	2.5	0.006
451	4.4	4	$3.10 \cdot 10^{-05}$
452	5.3	5	$2.90 \cdot 10^{-07}$
453	4.2	4	$3.10 \cdot 10^{-05}$
454	2.0	2	0.023
$p(n_{1,\min}, \dots, n_{5,\min})$			$3.85 \cdot 10^{-20}$
$P(n_1, \dots, n_5)$			$3.85 \cdot 10^{-17}$
$1/P(n_1, \dots, n_5)$			$2.60 \cdot 10^{16}$

washed with nitric acid from accumulating screens and from the walls of the experimental chamber and deposited on a powder of an aerosil. The peak was found only one time. Its chemical environment was composed of Ag, Ba, Nd, Y, Cd, Zr, and Sb. The spectrum containing it is given in Fig. 9.21. The nearest possible peak of “sums” was at a distance about 1 keV on the energy scale. The peak had a width of five channels, and its height reached 5.3σ at the maximum (see Table 9.14 and Fig. 9.22). It satisfies the 3σ criterion and could belong to peaks of the L- or M-series of chemical elements with numbers 141 to 143 and 260 to 267, respectively (see Table 9.9).

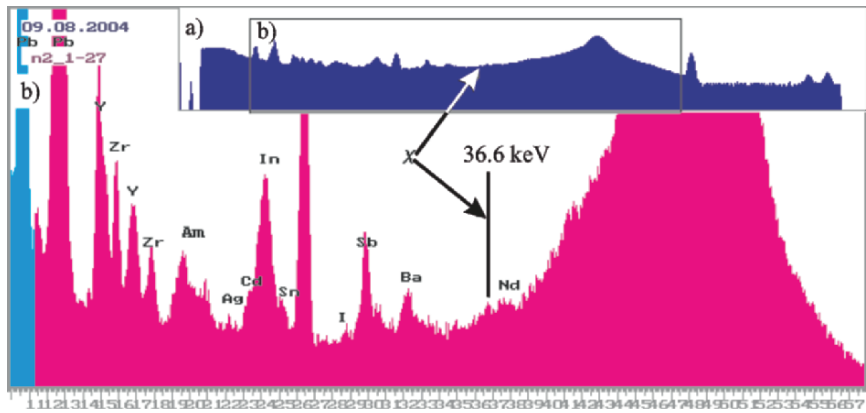


Fig. 9.21. X-ray spectrum containing the unidentifiable peak with an energy of 36.6 keV (a) and its fragment (b) presented on the logarithmic and linear scales, respectively.

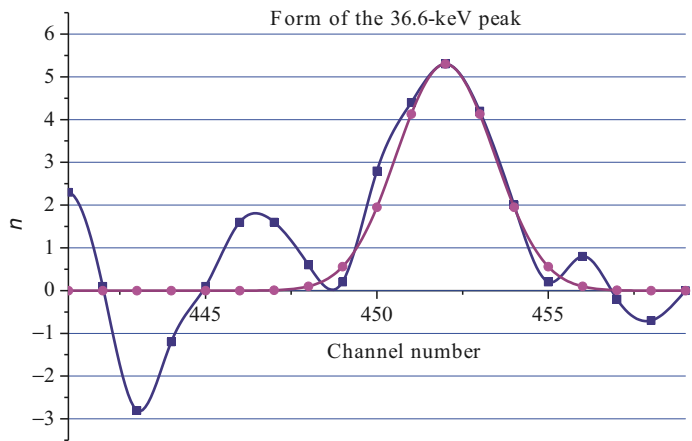


Fig. 9.22. Configuration of the unidentifiable peak with an energy of 36.6 keV as the dependence of the intensity in units of σ on the channel number of the X-ray spectrum (■) and its approximation with the normal distribution (●).

The unidentifiable peak with an energy of 38.1 keV, like the previous peak, was found on an installation for X-ray fluorescence analysis with the ^{241}Am isotope excitation source on specimen No. 1 (a deposit on aerosil). It was observed only one time. The spectrum containing it is given in Fig. 9.23. Its chemical environment was composed by Ag, Ba, Nd, La, Pb, In, Sn, Sb, Y, Zr, and Cd. There are no peaks of “sums” near it. However, it is partially overlapped by the Nd $K_{\alpha 1}$ peak from the low-energy side (see Fig. 9.24). This overlap occurs on the level that is below the half-height of the peaks,

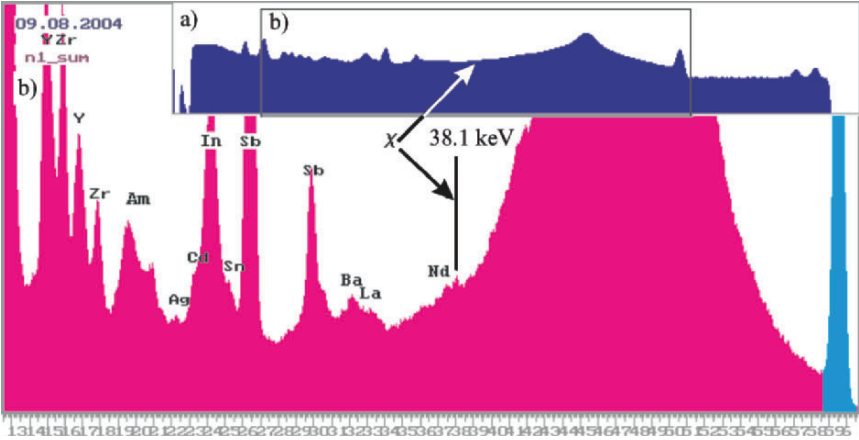


Fig. 9.23. X-ray spectrum containing the unidentifiable peak with an energy of 38.1 keV (a) and its fragment (b) presented on the logarithmic and linear scales, respectively.

Table 9.15. Probability of the appearance of a fluctuating peak with the configuration of the unidentifiable peak with an energy of 38.1 keV.

Channel number	$n_i = I_i/\sigma$	$n_{i,\min}$	$p_i(n_{i,\min})$
468	3.50	3.5	$2.3 \cdot 10^{-04}$
469	2.83	2.5	$6.0 \cdot 10^{-03}$
470	4.01	4	$3.1 \cdot 10^{-05}$
471	5.10	5	$2.9 \cdot 10^{-07}$
472	2.63	2.5	$6.0 \cdot 10^{-03}$
$p(n_{1,\min}, \dots, n_{5,\min})$			$7.4 \cdot 10^{-20}$
$P(n_1, \dots, n_5)$			$7.4 \cdot 10^{-17}$
$1/P(n_1, \dots, n_5)$			$1.34 \cdot 10^{16}$

which ensures their satisfactory resolution. In this case, the combination of the peaks is well approximated with a superposition of the corresponding normal distributions (see Fig. 9.24). The unidentifiable peak width is five channels, and it reaches the level of 5.1σ at the maximum (see Table 9.15). The peak is referred to the characteristic lines of the ^{241}Am isotope source and satisfies the 3σ criterion. Upon the extrapolation by the Moseley law, the peak falls into the L- and M-series of chemical elements with numbers 144 to 145 and 265 to 272, respectively (see Table 9.9).

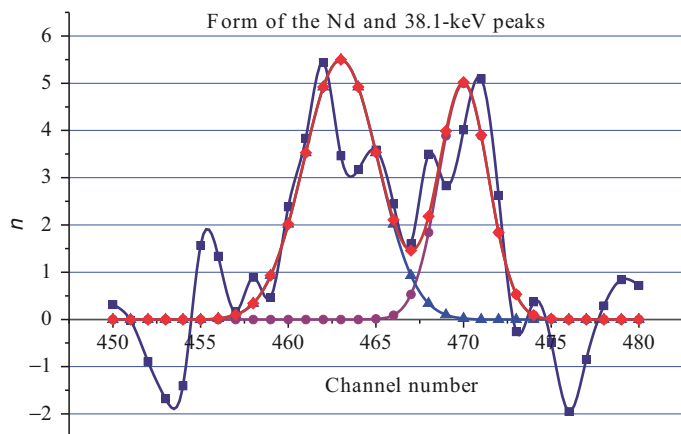


Fig. 9.24. Configuration of the unidentifiable peak with an energy of 38.1 keV as the dependence of the intensity in units of σ on the channel number of the X-ray spectrum (■) and its approximation with the normal distribution (●).

The unidentifiable peak with an energy of 51.95 keV was found on an installation for X-ray fluorescence analysis with the ^{241}Am isotope excitation source on the specimen, being a screen made of a polymeric film. Like two previous unidentifiable peaks, it was observed only one time. The spectrum containing the peak is given in Fig. 9.25. Its chemical environment is composed from Pb, Cu, Au, Fe, Ni, Sb, In, Sn, La, Ce, Cs, Pr, and Nd. The unidentifiable peak is located in the spectrum between the peaks of the elastically and inelastically scattered emission of the isotope source, where there are no peaks of losses and “sums” and the characteristic peaks of the source. The peak width was eight channels, and it reached the level of 4.4σ at the maximum (see Fig. 9.26). It satisfies the 3σ criterion (see Table 9.16) and could belong to characteristic peaks of the L-series or M-series of chemical elements with numbers 162 to 164 and 306 to 314, respectively (see Table 9.9).

All nine above-presented unidentifiable X-ray peaks were found during the study of the nucleosynthesis products over a period of a half-year. At present, we continue to search for them, and their number grows. It is necessary to notice that over all the time this study was carried out with use of specimens of two types. The first-type specimens were accumulating screens with target explosion products deposited on them in the experiments with an effective target damage, and the second-type specimens were accumulating screens with target substance deposited on them in the imitating experiments, in which there was no effective target damage.

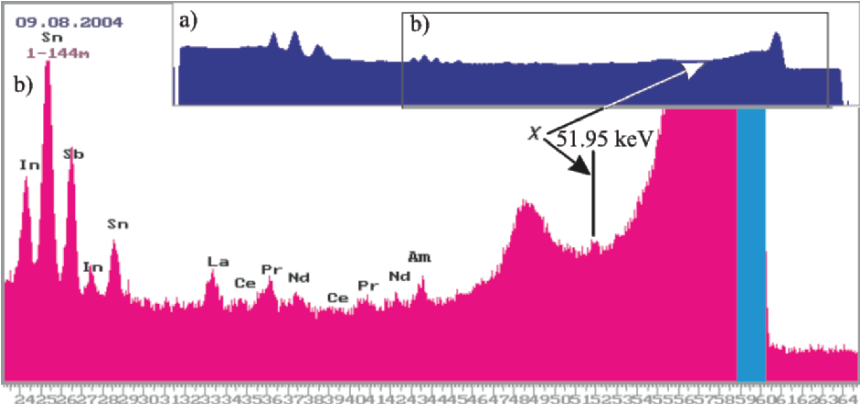


Fig. 9.25. X-ray spectrum containing the unidentifiable peak with an energy of 51.95 keV (a) and its fragment (b) presented on the logarithmic and linear scales, respectively.

Table 9.16. Probability of the appearance of a fluctuating peak with the configuration of the unidentifiable peak with an energy of 51.95 keV.

Channel number	$n_i = I_i/\sigma$	$n_{i,\min}$	$p_i(n_{i,\min})$
778	0.15	0	$1.00 \cdot 10^{00}$
779	4.40	4.00	$3.10 \cdot 10^{-05}$
780	3.23	3.00	$1.30 \cdot 10^{-03}$
781	2.17	2.00	$2.30 \cdot 10^{-02}$
782	3.06	3.00	$1.30 \cdot 10^{-03}$
783	3.38	3.00	$1.30 \cdot 10^{-03}$
784	2.71	2.00	$2.30 \cdot 10^{-02}$
785	1.22	1.00	$1.59 \cdot 10^{-01}$
$p(n_{1,\min}, \dots, n_{8,\min})$			$5.73 \cdot 10^{-18}$
$P(n_1, \dots, n_8)$			$5.73 \cdot 10^{-15}$
$1/P(n_1, \dots, n_8)$			$1.75 \cdot 10^{14}$

That is, there were nucleosynthesis products in the surface layer only of the first-type specimens. None of the X-ray microprobe operators who were engaged in the study knew about the existence of the accumulating screens “imitators”. However, all nine unidentifiable X-ray peaks were discovered by them only in the first-type specimens. This fact points out that the unidentifiable X-ray peaks are associated only with the nucleosynthesis products.

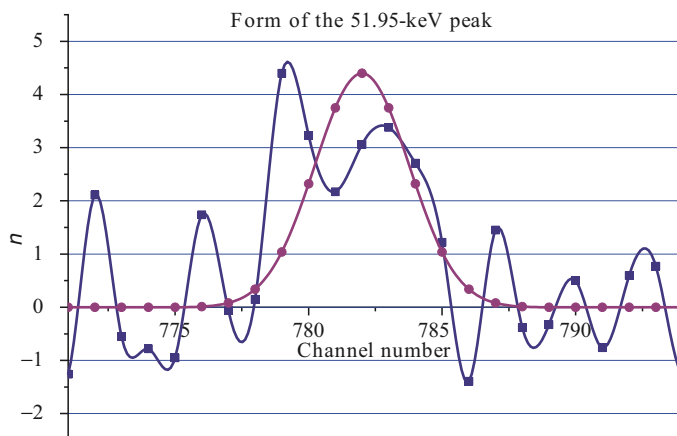


Fig. 9.26. Configuration of the unidentifiable peak with an energy of 51.95 keV as the dependence of the intensity in units of σ on the channel number of the X-ray spectrum (■) and its approximation with the normal distribution (●).

The detection of the unidentifiable X-ray peaks is an important argument in favor of that the nucleosynthesis products contain superheavy chemical elements, and the low intensity of the registered unidentifiable X-ray peaks in our judgment does not decrease the weight of this argument regarding their origin. Unidentifiable X-ray peaks exist without a doubt, and their low intensity is caused by the objective reasons comprehensively considered above.

9.2. Other Experimental Evidences for the Presence of Super-heavy Elements

Strictly saying, the discovery of superheavy chemical elements in target explosion products means, on the one hand, the observation of experimental events being evidences for the manifestation of their presence in products of the explosion and, on the other hand, properly their unambiguous identification as objects recognizable and existent, which assumes the determination of exact values of the masses and charges of their nuclei. It is necessary at once to recognize that the solution of the problem under study is quite far from completion at this time. This statement is related, in the first turn, to its last part, namely, to the identification of superheavy chemical elements. The complexity of this problem is conditioned by a number of objective reasons. First of all, the difficulty of its solution is connected now with the absence or the insufficient volume of knowledge about their chemical properties, the configurations of electron shells, the range of masses and

charges of nuclei, the amount and character of the distribution in products of the explosion, etc. The mentioned circumstances do not allow one to efficiently use, for their identification, traditional chemical methods, spectroscopic methods for various types of the characteristic emission generated upon a change of the state of the electron shells of atoms of chemical elements, and methods of mass-spectrometry separating the nuclides by the ratio of their masses to their charges. To the above, we mention the inapplicability of the known methods and scientific facilities to the study of such exotic and complex objects, as superheavy chemical elements.

As for the first part of the problem under discussion, the situation is different. Experimental events indicating the creation of superheavy chemical elements in explosive experiments are found upon the diagnostics of the plasma from a exploded target by detector methods and upon the study of the composition of products of the explosion which are deposited on the accumulating screens and walls of the chamber body upon the dispersion. Some of these events give direct evidences for the synthesis of superheavy chemical elements, and others indicate this circumstance indirectly. In the last case, we mean the situation where the results derived upon the studies cannot be consistently connected one with another without the assumption about the presence of superheavy chemical elements in products of the explosion. At present time, we have accumulated enough evidences to consider their existence as real. In this section, we present the most important from the mentioned evidences. In the cases where the information related to superheavy chemical elements was already presented in other sections of this collection, we will concern it exceptionally in the context of the comparative analysis concerning other analogous facts. But if the fact is given for the first time, we will consider it in more detail. In this case, we will not follow the chronological order upon the presentation of facts related to the problem under consideration.

9.2.1. Centralized Track Clusters with Anisotropic Distribution of Tracks

Let us consider the results of the registration of fast ions of the plasma expanding in the process of explosion-induced destruction of targets (see Sec. 6.5) with the help of track detectors (CR-39). In these investigations, as rare events, we registered the so-called centralized track clusters, consisting of the families of the tracks of α -particles and, possibly, other light nuclei with energies of up to 8 MeV and the common center of the dispersion. As usual, they contain from several tens to several hundreds of tracks. One of such track clusters, being unique and named “giant”, contained 276 tracks (see Sec. 6.5). All the mentioned track clusters are characterized by that

the tracks entering them have a strong directivity. The last means that the particles which formed these tracks in a detector had not only the common center of the dispersion, but possessed the momenta slightly different in the modulus and direction.

It is believed that the appearance of the nuclear decays registered by track detectors as centralized track clusters in target explosion products is referred to anomalous nucleophysical phenomena. By simulation, we have shown that similar track clusters cannot be created by fixed “hot” radioactive microparticles or dust particles upon any suitable exposure. In the last case, as distinct from our observations, one observes centralized track clusters, where the distribution of tracks is isotropic over azimuths. The analysis of the anisotropy of centralized track clusters shows that it is conditioned by the kinematics of a superheavy nuclear particle which has undergone instantaneous decay and moved with the velocity of the order of that characterizing the motion of track-forming particles. The centralized track clusters with anisotropic distribution of the tracks composing them are direct experimental evidences for the presence of superheavy chemical elements in target explosion products.

9.2.2. Instability of Unidentifiable X-Ray and Auger-Peaks under the Action of an Electron Probe

Let us now return to unidentifiable X-ray and Auger-peaks contained in the spectra registered on target explosion products (see Sec. 9.1.). Conjecturally, they are related to the characteristic X-ray and Auger emission of chemical elements positioned beyond the known part of the Periodic table, i.e., with superheavy chemical elements. This assumption is based, first of all, on the absence of any logical alternative in the explanation of the nature of the mentioned peaks. By concerning the problem in the context of superheavy chemical elements, we will focus our attention on one characteristic feature of all the discovered unidentifiable peaks: their instability under the action of an electron probe consisting in the “break-up” of peaks after their observation or their absence upon the repeated observations. At first glance, it seems that the indicated property is also characteristic of the peaks of the fluctuating nature appearing in spectra. However, unidentifiable peaks behave in a different manner. As distinct from fluctuating peaks, they are reproduced, but in the other places of registrations; and, after the “break-up” of peaks, the peaks of earlier absent chemical elements appear in the repeatedly observed spectra. In addition, the question about the fluctuating nature cannot be basically referred to unidentifiable Auger-peaks by virtue of that their amplitudes considerably exceeded the background fluctuations in the dominant majority of cases. It is also worth noting another feature in

the character of unidentifiable Auger-peaks. As distinct from X-ray peaks, we succeeded sometimes in observing them repeatedly at the same place, but their intensity decreased remarkably in these cases.

The described behavior of unidentifiable peaks testified to the favor of the correctness of the assumption that they belong to the characteristic peaks of superheavy chemical elements and allows us to think that it is nothing but a manifestation of the decay of superheavy chemical elements under the perturbing action of an electron probe. This was indicated by, first of all, the very “decay” of unidentifiable peaks, the appearance of new chemical elements as its result, and their clearly more tolerant reaction to the action of the electron beam of an Auger-microprobe as compared to the X-ray microprobe. The last circumstance was related, we believe, to a smaller accelerating voltage of the electron beam of an Auger-microprobe and, hence, with a smaller perturbing action as compared to that of an X-ray microprobe. We recall that their accelerating voltages were 10 and 35 kV, respectively.

With the purpose of verifying the above assumption, we planned and performed the testing experiments with the using detectors to observe possible nuclear particles emitted in the process of decay of nuclei which are hypothetically present in products of the explosions of superheavy chemical elements, the decay being initiated by the beams of low-energy particles. Below, we give the comments to these experiments and note here that their result was positive. This indicates the correctness of the assumption that unidentifiable X-ray and Auger-peaks belong to the characteristic peaks of superheavy chemical elements and puts them in the series of experimental facts directly testifying to the presence of superheavy chemical elements in target explosion products.

9.2.3. Initialization of High-Energy Nuclear Particle Emission by Low-Energy Beam Irradiation

The above-mentioned experiment was carried out in two versions (see Sec. 10). In one of them, the explosion products of lead targets, which were deposited on copper accumulating screens, were irradiated by the beam of ions O_2^+ with an energy of about 12.5 kV and a beam current of 1 to 5 μA . In the second version, the irradiation was realized by a laser emission with a wavelength of 1.06 μm . In this case, the emission energy in a pulse was about 50 mJ, the pulse duration was 14 ns, and the emission power density on the screen was $2 \times 10^9 W/cm^2$. As detecting tools, we used a Si detector of α -particles and a track detector CR-39, respectively. In both experiments, we observed the emission of high-energy nuclear particles, which was

stimulated by the low-energy irradiation, with intensity significantly exceeding that of the natural α -background. In the case of the irradiation by ions O_2^+ where the observation was made in real time, we discovered two scenarios of the running of the process of stimulated emission. In the first scenario, the emission of high-energy nuclear particles was looked as a flow of random separated and relatively rare events lasting from 1 to 2 hours. For the mentioned time interval, we usually registered several tens of such particles. For them, the product of energy by charge was several tens of times more than that of background α -particles, whose counting rate was about 0.5 particle/h. In the second scenario, the emission was looked as a flow of random relatively rare series of events with their duration of about several seconds each. In this case, during one series, the emission of several hundreds of particles was usually observed, and the product of energy by charge was only by several times more than that of background α -particles.

It is quite obvious that the observation of nuclear particles with the observed energies and frequencies, which considerably exceed the mentioned characteristics of α -particles of the natural background, give indication that their sources are some sort of nuclear reactions running in the studied specimens under the action of irradiation. By virtue of that the nuclei of ordinary chemical elements cannot enter nuclear reactions under the action of low-energy particles or laser emission, the results of the executed experiments do imply that the target explosion products contain superheavy chemical elements. Moreover, the same results indicate the fact that their nuclei have a property quite unusual by the modern physical ideas: the possibility to enter into nuclear reactions due to the action of low-energy perturbations. As for the type of the mentioned reactions, we do not have sufficient foundations to consider that they are related exclusively to the decay of the metastable nuclei of superheavy chemical elements, which appears as the most natural scenario to explain situation under consideration. Apparently, there also exist the nuclear reactions of a step-by-step synthesis of heavier nuclei running due to the interaction of an initial "maternal" superheavy nucleus with surrounding nuclei of the substance of an accumulating screen, which is accompanied by the emission of registered light high-energy nuclear particles. The indicated type of nuclear reactions corresponds, most probably, to the second scenario of the running of the process of stimulated emission of high-energy nuclear particles. Below, we present some experimental facts indirectly supporting the above assumption.

9.2.4. Nonfulfillment of the Energy Balance in the Running Nuclear Transformations from the Composition of Nucleosynthesis Products

We now consider a number of experimental results indirectly indicating the presence of superheavy chemical elements in target explosion products. Let us begin from the results of the determination of the element composition of target explosion products by glow-discharge mass-spectrometry (see Sec. 8.2.1). The main results are the following: upon the explosive destruction of a target, about 3.5×10^{18} to 8.0×10^{18} atoms of new chemical elements, which contain about 1.7×10^{20} to 4.6×10^{20} nucleons on the average, are generated. In this case, the share of nucleosynthesis products in target explosion products is in the scope from 10% to 26.5%, and the mean number of nucleons per one atom of the registered part of the regenerated target substance varies approximately from 33 to 58. These data were derived in the experiments with targets made of Al, Fe, Ag, and Pb. The basic contradiction revealed upon the analysis of the data is the nonfulfillment of the energy balance in the running nuclear transformations.

Let us begin from the experiments with iron targets. Since the specific binding energy of a nucleus of the target substance exceeds the mean specific binding energy of nuclei of the reaction products, we are faced with a deficit of the total energy yield of the reaction of about 3.1×10^3 kJ. This deficit cannot be compensated because the energy store of the setup is at most 2.5 to 3.0 kJ and the energy of the shock action of an electron beam on a target is in the range from 100 to 300 J by the results of systematic electric measurements. The last circumstance indicates the impossibility for the considered nuclear transformations to have occurred by “standard” nuclear physics. However the experimental evidence has shown that they are occurring! As for the experiments with Al, Ag, and Pb targets, we should have a positive huge energy yield from the reactions in view of the observed composition of nucleosynthesis products. It is in the scope of 4.5 to 31 MJ, which corresponds the energy release of 1.1 to 7.4 kg of TNT upon the explosion (the specific heat of an explosion of TNT equals 4.2 MJ/kg [see Refs. 78, 82]). It is obvious that the release of such amounts of energy cannot occur without consequences for the experimental setup and go unnoticed.

In order to understand the reasons for the mentioned contradictions, it is necessary, first of all, to consider the scheme of calculations underlying the estimations of the energy balance of nuclear transformations running upon the explosive destruction of targets. We will briefly characterize the situation as follows. First, on the basis of direct measurements, we determined the number of atoms of various chemical elements generated as a result of the nucleosynthesis and calculated the number of nucleons contained in

them. Basing on the assumption that the number of nucleons is conserved in the transformations under consideration, we calculated the number of atoms of a target substance, which participated in nuclear reactions. Then the energy yield of reactions was determined as the difference of the total binding energies of nucleons in regenerated atoms, on the one hand, and atoms of the target substance which entered the nucleosynthesis, on the other hand.

While discussing the presented scheme of calculations, it is necessary to give attention to the circumstance that the discrepancies in the energy balance of reactions cannot be a consequence of even improbably great errors in the determination of the observed composition and, respectively, the amount of nucleosynthesis products. For example, by assuming that the amount of nucleosynthesis products in the experiments with aluminum targets is overestimated by one order of magnitude, we get, nevertheless, the energy yield in one experiment to be equivalent to that upon the explosion of 740 g of TNT. In other words, the reasons for the seeming contradictions must lie in other place. In a formal way, the huge discrepancies in the calculations of energies can be related to two factors: on the one hand, to a large differences in the specific binding energies of nucleons in registered products of the nucleosynthesis and in the initial target substance and, on the other hand, to the use of the assumption that the number of nucleons is conserved in the occurring nuclear transformations. Upon the analysis of the role of the mentioned factors in the origin of discrepancies in the energy balance, it becomes obvious at once that the assertion about the conservation of the number of nucleons should not be questioned. But it should be correctly used. However, the unique way to remove these contradictions is the search for the reason of the great difference in the specific binding energies of nucleons in the registered nucleosynthesis products and the initial target substance. In other words, we may say that we need to seek the means to diminish this difference.

The unique way to diminish the difference in the specific binding energies under consideration is to assume that the products of nucleosynthesis contain, besides the observed part, a part which cannot be discovered within the applied methods of determination of the composition and represents superheavy chemical elements. It is obvious that, by varying the specific binding energy of nucleons in the nuclei of this undetected part of nucleosynthesis products and by using the assertion about the conservation of the number of nucleons, it is possible to easily ensure the observance of the energy balance in all the transformations under study. For example, for the energy balance to be satisfied, it is necessary that the binding energy of nucleons in the unobserved part of nucleosynthesis products in the experiments with iron targets be more than 8.787 MeV. In the experiments with Pb, Ag, and Al targets, it should be

less: 7.870, 8.551, and 8.332 MeV, respectively. In other words, the interpretation of results of the determination of the composition of nucleosynthesis products by glow-discharge mass-spectrometry requires the presence of the unregistered part of nucleosynthesis products, which indicates indirectly the presence of superheavy chemical elements in them.

9.2.5. Divergence of the Amount of a Target Matter with its Observed Amount on the Accumulating Screens

Solving the contradictions arising upon the calculation of the energy balance of the considered nuclear transformations, we drew conclusion, by means of speculations, that the products of nucleosynthesis deposited on the accumulating screens include the part which is unregistered within the applied methods of determination of a composition. The indicated conclusion could be set on a stronger basis, if we would find, in addition to a logical analysis, some experimental facts confirming it. The earlier described scheme of calculations of the energy balance yields that such unique argument, which is based on the experimental measurements and indicates the presence of unregistered superheavy chemical elements in the nucleosynthesis products, could be the establishment of the fact that the amount of a target substance entering into nuclear reactions exceeds its observed amount in products of the mentioned reactions.

To determine the amount of a substance entering into nuclear reactions upon the explosions of targets, we can use our results from experiments on the neutralization of radioactivity (see Sec. 7). The method of its determination is based on the following suppositions. The many-year experience of the execution of explosive experiments showed that, irrespective of the initial target substance, the products of nuclear regeneration contain always only stable nuclides in a significant amount. According to the above-written, in the cases where a radioactive substance was taken as an initial target substance, its radioactivity drops by a value corresponding to the share of the regenerated substance as a result of the explosion. Thus, to correctly evaluate the amount of the substance entering the nuclear reactions, it is necessary to use a one-piece uniformly radioactive tip of a target with the known mass and initial radioactivity. In order that the regeneration zone does not go beyond the limits of the radioactive tip (otherwise, the evaluation will be underestimated), its length should considerably exceed its diameter. It is desirable that their ratio be at least 5 to 1. Upon the observance of these conditions, the neutralization degree of radioactivity will correspond, as a result of the explosion, to the share of the substance mass entering the reactions in the starting mass of the radioactive tip.

We evaluated the amount of the substance entering the nuclear reactions by using the targets with one-piece cobalt tips containing uniformly distributed radioactive nuclide ^{60}Co . The ratio of their length to a diameter was about 12. The change of radioactivity was registered by two lines of γ -emission of ^{60}Co with energies of 1173.2 and 1332.5 keV. The neutralization degree averaged by the results of several experiments was about 14.4%. The data derived yields that the number of Co atoms entered the nuclear reactions was at the level of 1.6×10^{19} per one target 0.5 mm in diameter, and they contained about 9.4×10^{20} nucleons. It is easy to see from the comparison of the derived data with the data of glow-discharge mass-spectrometry on the amount of nucleosynthesis products that the number of nucleons contained in the substance which entered the nuclear reactions exceeds their number in the registered part of the reaction products by 2 to 5 times. This means that the discovered excess of nucleons appeared as a result of the running nuclear reactions is bound in the unregistered part of nucleosynthesis products, most probably in the nuclei of superheavy chemical elements.

We now consider the results of the determination of the composition of a substance dispersed after the explosion of a target by the optical spectra of emitting ions contained in a plasma bunch expanding from the center of the explosion (see Sec. 4). The most important results are the following: a plasma bunch includes about 1.8×10^{17} to 6.1×10^{17} emitting ions which contain about 0.58×10^{19} to 4.3×10^{19} nucleons on the average. In this case for the dominant share of experiments, the average number of nucleons per ion of the plasma bunch is in the interval 42 to 71. These data were derived on a large sampling of experiments with targets made of Al, Cu, and Pb, in which Al, Cu, and Ta served as the substance of accumulating screens.

Upon the comparison of the presented data with analogous ones derived by glow-discharge mass-spectrometry, it is worth noting, first of all, the fact that the number of nucleons in the regenerated atoms registered on the accumulating screens exceeds the number of nucleons in the ions contained in a plasma bunch by more than one order of magnitude. This fact cannot be explained by the measurement errors. The determination of the composition of plasma bunches and of nucleosynthesis products on the accumulating screens was repeated many times, and the results were steadily reproduced. Moreover, the derived difference cannot be referred to a higher sensitivity of the method of glow-discharge mass-spectrometry. Indeed, optical methods cannot observe the ions of chemical elements with a small content in a plasma bunch. However, this circumstance cannot lead to such a difference since, in both cases, at most 10 chemical elements cover 90% of the amount of observed atoms.

In other words, the derived difference of the results allows us to conclude that a plasma bunch includes also the unregistered part consisting of superheavy chemical elements. Generally saying, this conclusion is obvious. Indeed, if superheavy elements are contained in the products of nucleosynthesis on the accumulating screens, they must be present also in a plasma bunch, since their synthesis occurs in it on a certain stage of its evolution. The marked difference observed between the results in the plasma bunch as compared to the results on the accumulating screens only indicates the circumstance that the share of superheavy chemical elements in the products of the regenerated target substance is significantly greater during the stage of the expansion of a plasma bunch than in the state of precipitation on an accumulating screen.

Of interest is also another difference of the compared results. The mean number of nucleons per ion in a plasma bunch exceeds significantly an analogous number for the atoms contained in the products of nucleosynthesis deposited on the accumulating screens. This fact could be conditioned by the decay of a part of superheavy chemical elements synthesized in a plasma bunch on their collisions with atoms of the substance of an accumulating screen which is accompanied by the emission of light nuclei. The proposed scenario explains a decrease in the share of superheavy chemical elements in the nucleosynthesis products deposited on the accumulating screens and the shift of the nucleosynthesis products composition to the side of lighter chemical elements. Its reality is supported by the above-discussed results of experiments studying the decay of superheavy chemical elements which is stimulated by low-energy perturbations. Indeed, upon the collisions with frame ions of the matrix of an accumulating screen, superheavy atoms scattered by them upon the deposition undergo the perturbations not lesser than those in the case of their irradiation by a beam of ions O_2^+ with an energy of about 12.5 keV. In this case, it is obvious that the density of frame ions of the matrix exceeds significantly the density of ions O_2^+ in the irradiated beam. The last circumstance will promote the more efficient decay of superheavy chemical elements upon their deposition on accumulating screens as compared to that which was registered upon the irradiation of accumulating screens by an ion beam.

9.2.6. Anomalies in the Isotope Composition of the Material of Accumulating Screens

To the facts indirectly testifying to the presence of superheavy chemical elements in target explosion products deposited on the accumulating screens, we can refer, apparently, the anomalies of the isotope composition of a basic chemical element of the material of accumulating screens, being registered

after explosive experiments. As an example of such an anomaly, we have presented the deviation of the content of nuclide ^{65}Cu by 4.9% from its natural abundance in Sec. 8.1.2. It was observed on a specimen of the “sandwich” type which was made of copper accumulating screens of 0.5 mm in thickness. By discussing the presented result we note, first of all, that the mentioned anomaly exceeds the measurement error almost by one order of magnitude. At the same time, the explanation of the origin of this anomaly meets with difficulties. Indeed, the mean mass loss of a target upon the explosion was 2.6 mg in the experiments under discussion, and the mass of an accumulating screen was about 790 mg. Even if we assume that all lead of the target transferred on the accumulating screen upon its explosion is regenerated in one of the nuclides of Cu, its amount in this extreme case will still be too small in order to change so much the isotope composition of the massive copper screen due to any mixing phenomenon.

The presented arguments yield that any interpretation of this fact must be related only to the nuclear regeneration of the very substance of an accumulating screen as a result of the explosion of a target. In our opinion, the most probable scenario of such a regeneration can involve the above-mentioned nuclear reactions of a step-by-step synthesis of heavier nuclei by means of the interaction of a primary superheavy nucleus, which got into the screen substance with target explosion products, with screen substance nuclei surrounding it. Only reactions of such a type can provide the interaction of one superheavy nucleus with a large number of nuclei of the basic chemical element of the substance of an accumulating screen and can induce a considerable change in its isotope composition. Possibly, just these reactions were registered in the case of the second scenario realized in the experiments on the stimulation of the nuclear transformations of superheavy chemical elements by a low-energy beam of ions. It is obvious that the emission of a huge number of high-energy nuclear particles as a series for a short-time interval looks like a more natural manifestation of such a synthesis, than the emission of a small number of particles as individual rare events.

9.2.7. Qualitative Differences of the Observed Compositions of a Plasma Bunch and Nucleosynthesis Products Deposited on Accumulating Screens

So far, discussing the results of the determination of the compositions of the target explosion products deposited on accumulating screens and a plasma bunch, we have concentrated mainly on their quantitative aspect: we have compared the numbers of atoms and nucleons contained in them which were observed in various measuring experiments. We now focus our attention

on that from which ions a plasma bunch is composed and which atoms of chemical elements belong to the explosion products deposited on the accumulating screens. In other words, we will perform the qualitative comparison of the results of the determination of the composition of objects studied with various methods. It is more convenient and most natural to make the comparison by using the tools of statistical analysis, than by any other means.

We are interested, first of all, in the degree of coincidence or difference of the character of a distribution of chemical elements in the measuring experiments on the registration of the composition of the studied objects, i.e., in their correlation. The correlation analysis of the data on the composition is needed in order to answer a number of questions extremely important for the comprehension of how the process of nucleosynthesis upon the target explosions occurs, beginning from the stage of expansion of a plasma bunch and completing the stage of its condensation on the accumulating screen. This analysis is important for the clarification of the character of the dependence of the composition of nucleosynthesis products in the process of evolution on the compositions of the initial substance of a target and an accumulating screen. Upon the consideration of these questions, it is convenient to group all the analyzed results by their relation to the type of a studied object. To the first and second groups, we refer the data on the determination of the composition, respectively, of the ion component of a plasma bunch by optical-spectral methods and of the surface layer of nucleosynthesis products deposited on accumulating screens by X-ray electron probe microanalysis and Auger-electron spectroscopy. And the third group includes the data derived for the whole volume of accumulating screens together with deposited explosion products by glow-discharge mass-spectrometry.

In mathematical statistics, the pairwise connection of random variates of a system is characterized by correlation moments (see Refs. 144–147)

$$K_{ij} = M[(x_i - m_i)(x_j - m_j)], \quad i \neq j, \quad (9.8)$$

where x_i and m_i are values taken by the i -th random quantity and its expectation, respectively. In a particular case at $i = j$, the correlation moment Eq. 9.8 is nothing but the variance of the i -th random value:

$$K_{ii} = D_i. \quad (9.9)$$

It is convenient to position all correlation moments and variances in the form of a rectangular table called the correlation matrix of a considered system of random quantities:

$$\begin{pmatrix} K_{11} & K_{12} & \dots & K_{1n} \\ K_{21} & K_{22} & \dots & K_{2n} \\ \dots & \dots & \dots & \dots \\ K_{n1} & K_{n2} & \dots & K_{nn} \end{pmatrix}. \quad (9.10)$$

It is obvious that not all the elements of the correlation matrix are different. From the definition of correlation moment, it becomes clear that $K_{ij} = K_{ji}$, i.e., the elements of the correlation matrix positioned symmetrically relative to the main diagonal are equal. In this connection, one fills usually not the whole correlation matrix, but only its upper half by counting from the main diagonal:

$$\begin{pmatrix} K_{11} & K_{12} & \dots & K_{1n} \\ & K_{22} & \dots & K_{2n} \\ & & \dots & \dots \\ & & & K_{nn} \end{pmatrix}. \quad (9.11)$$

The correlation matrix composed from elements K_{ij} is often denoted by the symbol $\|K_{ij}\|$ for brevity.

In the case where the considered random values are not correlated, all the elements of a correlation matrix, except for diagonal ones, are equal to zero. On the contrary, if nondiagonal elements of the matrix are different from zero, the system of random values is correlated. For the sake of clearness of the assertion about the correlated nature of random values irrespective to their variance, the correlation matrix $\|K_{ij}\|$ is substituted by a normalized correlation matrix $\|r_{ij}\|$ composed from correlation coefficients, rather than from correlation moments:

$$r_{ij} = K_{ij}/\sigma_i\sigma_j, \quad (9.12)$$

where

$$\sigma_j = \sqrt{D_j} \quad \text{and} \quad \sigma_i = \sqrt{D_i}. \quad (9.13)$$

All the diagonal elements of this matrix are equal, naturally, to 1. Thus, a normalized correlation matrix takes the form

$$\|r_{ij}\| = \begin{pmatrix} 1 & r_{12} & r_{13} & \dots & r_{1n} \\ & 1 & r_{23} & \dots & r_{2n} \\ & & 1 & \dots & r_{3n} \\ & & & \dots & \dots \\ & & & & 1 \end{pmatrix}. \quad (9.14)$$

In Table 9.17, we give the results of the pairwise comparison of the distributions of chemical elements in the compositions of the ion component of a plasma bunch which are determined by optical-spectral analysis in the experiments with different target/screen compositions. In essence, this table is presented in the form of the above-described normalized correlation matrix. In the cases where the word “screen” is absent in the designations of

Table 9.17. Correlation coefficients of the distributions of chemical elements in the compositions of a plasma bunch in experiments with different target/screen compositions.

	Pb/Cu	Pb/Ta	Pb/Al	Pb/	Cu/Cu	Cu/Cu	Cu/Cu	Cu/	Cu/	Al/Cu	Al/Ta	$M_i^\dagger$	$M^\ddagger$
Pb/Cu*	1	0.93	0.79	0.86	0.91	0.90	0.77	0.89	0.75	0.95	0.95	0.87	
Pb/Ta		1	0.97	0.81	0.90	0.85	0.85	0.82	0.82	0.96	0.52	0.84	
Pb/Al			1	0.84	0.80	0.76	0.50	0.75	0.75	0.84	0.52	0.75	
Pb/				1	0.90	0.88	0.85	0.87	0.79	0.94	0.84	0.86	
Cu/Cu					1	0.98	0.88	0.77	0.86	0.87	0.90	0.86	
Cu/Cu						1	0.87	0.81	0.84	0.83	0.86	0.85	0.81
Cu/Cu							1	0.84	0.91	0.41	0.72	0.76	
Cu/								1	0.65	0.43	0.46	0.75	
Cu/									1	0.79	0.79	0.78	
Al/Cu										1	0.97	0.74	
Al/Ta											1	0.75	

* target/screen
† mean over a composition
‡ mean over a group

a composition in the texts to rows and columns of the matrix, the explosive experiments were realized without an accumulating screen. In the penultimate column, we give the mean correlation coefficients for the compositions derived in the experiments indicated in the relevant rows. Finally, the last column contains the mean correlation coefficients of compositions of the whole group.

We now consider properly the data presented in Table 9.17. As for them, we give only one comment: all, without exception, correlation coefficients are quite close to 1. This means that the compositions of the ion component of plasma bunches slightly differ one from another in all the performed explosive experiments. In view of this fact and the circumstance that we used a wide collection of various substances in the production of targets and accumulating screens, we may conclude that the composition of a plasma bunch is practically independent of the substances of a target and a screen. This independence is understandable, because the target explosion products did not interact with an accumulating screen prior to their deposition. But the absence of the connection between the composition of a

Table 9.18. Correlation coefficients of the distributions of chemical elements in the compositions of the surface layers of screens in experiments with different target/screen compositions.

	Ag/Cu	Al/Cu	Ni/Cu	Pb/Cu	Pd/Cu	Ta/Cu	V/Cu	Cu/Al	Cu/Cu	Cu/V	Cu/Zr	$M_i^\dagger$	$M^\ddagger$
Ag/Cu*	1	0.89	0.89	0.83	0.95	0.91	0.51	0.61	0.82	0.77	0.04	0.72	
Al/Cu		1	0.96	0.87	0.82	0.67	0.70	0.72	0.82	0.74	0.13	0.73	
Ni/Cu			1	0.82	0.81	0.69	0.73	0.77	0.91	0.73	0.26	0.76	
Pb/Cu				1	0.70	0.38	0.39	0.62	0.57	0.54	0.19	0.59	
Pd/Cu					1	0.95	0.45	0.53	0.74	0.71	-0.08	0.66	
Ta/Cu						1	0.32	0.41	0.73	0.72	0.00	0.58	0.63
V/Cu							1	0.46	0.78	0.76	0.48	0.56	
Cu/Al								1	0.67	0.53	0.59	0.59	
Cu/Cu									1	0.80	0.53	0.74	
Cu/V										1	0.39	0.67	
Cu/Zr											1	0.25	

* target/screen

† mean over a composition

‡ mean over a group

plasma bunch and the substances of a target is, indeed, a rather unexpected observation.

Table 9.18 has the same structure as the previous table. The former is a normalized correlation matrix of the distributions of chemical elements in the compositions of the surface layers of nucleosynthesis products on the accumulating screens which are registered by X-ray electron probe microanalysis and Auger-electron spectroscopy in experiments with different target/screen compositions. By analyzing the data presented in Table 9.18, it is worth noting that the situation is similar to that in the previous case. Here, the correlation coefficients, while being somewhat lower than those in the previous case, are relatively high. For example, the mean group correlation coefficient of the compositions equals 0.63 as compared to the result in the first case of 0.81. Attention should be drawn to the fact that the composition Cu/Zr is not correlated with other target/screen compositions. However, we should hardly interpret it as a strong dependence of the composition of the surface layer of nucleosynthesis products on the substance of an accumulating screen, such a dependence being not explicitly observed in all the cases.

At the same time, we have not found any satisfactory interpretation of this fact.

Thus, we may draw the following conclusion: the composition of a layer of nucleosynthesis products on the accumulating screens has no pronounced dependence on the substances of a target and an accumulating screen. It is obvious that this circumstance is conditioned, first of all, by the constancy of the composition of a plasma bunch. Indeed, in the surface layers of accumulating screens, we registered mostly the condensate of a plasma bunch which is distinguished, as was indicated above, by the stability of its composition. At the same time, it is necessary to note that the fact of a lesser correlation of the compositions of the surface layers of nucleosynthesis products is related, apparently, to a certain effect of a substance of screens on them.

Finally, we move to the comparison of the compositions of nucleosynthesis products observed by glow-discharge mass-spectrometry in the volume of the whole screen in experiments with various target/screen compositions. These data are also presented in the form of a normalized correlation matrix in Table 9.19. Upon the analysis of these data, we pay attention to, first of all, the fact that the indicated matrix has a clearly pronounced block structure. These blocks are singled out in Table 9.19 by lines. The upper left and lower right diagonal blocks demonstrate the quite high values of correlation coefficients. For them, the mean correlation coefficients are 0.67 and 0.81, respectively. In these blocks, we can compare the compositions

Table 9.19. Correlation coefficients of the distributions of chemical elements in the compositions of nucleosynthesis products of the whole volume of screens in experiments with different target/screen compositions.

	Al/Cu	Fe/Cu	Cu/Cu	Zn/Cu	Ag/Cu	Pb/Cu	Fe/Nb ₁	Fe/Nb ₂	Fe/Nb ₃	$M_1^\dagger$	$M_2^\dagger$	$M_3^\ddagger$	$M_4^\ddagger$
Al/Cu*	1	0.98	0.45	0.87	0.90	0.99	0.13	0.27	0.23	0.60	0.84		
Fe/Cu		1	0.51	0.48	0.93	0.96	0.46	0.41	0.43	0.65	0.77		
Cu/Cu			1	0.30	0.42	0.41	0.30	0.76	0.53	0.46	0.42		
Zn/Cu				1	0.30	0.69	0.19	0.48	0.39	0.46	0.53	0.67	
Ag/Cu					1	0.89	0.09	0.21	0.18	0.49	0.69		
Pb/Cu						1	0.09	0.17	0.15	0.54	0.79		
Fe/Nb ₁							1	0.61	0.90	0.20	0.76	0.48	
Fe/Nb ₂								1	0.90	0.48	0.76	0.81	
Fe/Nb ₃									1	0.46	0.90		

* target/screen
† mean over a composition (1 – for a matrix, 2 – for a block)
‡ mean over a group (3 – for a whole group, 4 – for a subgroup)

derived in experiments with the same substance of an accumulating screen: in the first and second blocks, the initial substances of screens were, respectively, copper and niobium. The upper right block shows the very low values of correlation coefficients. In it, we can compare the compositions of nucleosynthesis products observed in our experiments with various initial substances of accumulating screens.

Thus, the above-presented discussion yields that the initial substance of an accumulating screen has a considerable influence on the composition of nucleosynthesis products registered in the volume of a whole screen, as distinct from the target substance. This assertion should be commented. The established dependence cannot be interpreted, in principle, as a manifestation of an admixture contained in the initial substance of an accumulating screen. The mentioned interpretation is impossible due to two reasons. First, the observed amount of nucleosynthesis products in the volume of a whole screen exceeded, as a rule, the amount of an admixture of the initial substance of a screen almost by two orders of magnitude. Second, upon the calculation of the amount of nucleosynthesis products, we always subtracted the amount of an admixture which was contained in both the initial accumulating screen and the transferred target substance from the registered amount of “foreign” chemical elements contained in the accumulating screen after the experiment. In other words, the discussed dependence can be explained only by the interaction of the deposited products of nucleosynthesis with the initial substance of the accumulating screen. In this case, it is obvious that the indicated interaction must include nuclear transformations in order to induce the observed changes of the composition of the deposited products of nucleosynthesis. Such nuclear transformations in an accumulating screen can be ensured only by the presence of superheavy chemical elements in deposited target explosion products, since the nuclei of ordinary chemical elements do not enter nuclear reactions under such conditions.

Up to now, we have compared the compositions of nucleosynthesis products in experiments with different target/screen compositions in the scope of a single studied object: a plasma bunch, the surface layer, and the whole volume of an accumulating screen. This allowed us to reveal the influence of the initial substances of a target and an accumulating screen on the composition of nucleosynthesis products on different stages. In order to more completely represent the character of the evolution of a composition of nucleosynthesis products, it is useful to compare their compositions on the mentioned stages. Such a comparison can be carried out by using the data presented in Table 9.20. The diagonal elements in this table are the mean intragroup correlation coefficients of the compositions in a plasma bunch, in the surface layer of nucleosynthesis products on an accumulating screen,

Table 9.20. Mean intra- and intergroup correlation coefficients of the distributions of chemical elements in the compositions of nucleosynthesis products registered in a plasma bunch, the surface layer, and the whole volume of screens in experiments with different target/screen compositions.

Object of analysis	Plasma bunch	Surface layer of nucleosynthesis products on the accumulating screen	Nucleosynthesis products in the volume of an accumulating screen
Plasma bunch	0.81	0.64	0.20
Surface layer of nucleosynthesis products on an accumulating screen		0.63	0.36
Nucleosynthesis products in the volume of an accumulating screen			0.48

and in the products of nucleosynthesis in the volume of an accumulating screen, respectively. It is easy to see that the indicated correlation coefficients in the row r_{11} , r_{22} , r_{33} decrease from 0.81 to 0.48 with increase in the index. This means that, upon the transition from one stage to the other, the composition of nucleosynthesis products becomes increasingly diverse. At the same time, it is necessary to note that the correlation of the compositions of nucleosynthesis products in the volume of an accumulating screen remains quite considerable as before, $r_{33} = 0.48$, even on the last considered stage.

We now consider the nondiagonal elements of the matrix under study (Table 9.20), i.e., the intergroup correlation coefficients. A high value of the coefficient r_{12} equal to 0.64 give evidence that the observed compositions of a plasma bunch and nucleosynthesis products in the surface layer of an accumulating screen differ slightly. This confirms the above-presented assumption that the upper layer of nucleosynthesis products on the accumulating screen is the condensate of a plasma bunch in an almost unchanged form. The intergroup coefficient r_{13} characterizes the correlation of the observed compositions of a plasma bunch and nucleosynthesis products in the volume

of an accumulating screen. Its value is rather low and equals 0.20. It would be, most probably, lower if a wider collection of substances would be used in the production of accumulating screens in our experiments. This means that the last compositions have almost nothing in common. Thus, a sharp change of the composition of nucleosynthesis products occurs, apparently, after the penetration of superheavy chemical elements into the volume of a screen and due to their interaction with its substance. We believe that this interaction leads to the creation of an amount of new nucleosynthesis products so large that the unmodified or observed part of the initial nucleosynthesis products deposited on the surface of an accumulating screen “sinks” in them, and the connection with the prehistory in the evolution of the composition of nucleosynthesis products is basically lost.

9.2.8. Layers of Anomalous Enrichments in the Accumulating Screens

Thus, both the comparative analysis of the compositions of a plasma bunch and nucleosynthesis products in the surface layer and in the volume of accumulating screens and the observed anomalies of the isotope composition of a substance of screens arising due to the explosions of targets lead us to the conclusion that superheavy chemical elements penetrate into the volume of an accumulating screen upon the condensation of a plasma bunch. In this case, they undergo nuclear transformations there by decaying and/or by entering the reactions of a step-by-step synthesis of heavier nuclei by means of the interaction with nuclei of a screen substance surrounding them. We mention one more independent experimental evidence confirming both the described scenario of the penetration of superheavy chemical elements into the volume of an accumulating screen and the character of their behavior in it. That evidence is the data presented by the study of the distribution of nucleosynthesis products over the depth of accumulating screens (see Sec. 11.2.6) by secondary-ion mass-spectrometry (IMS 4f, CAMECA, France).

It was expected that the maximum amount of the condensate of a plasma bunch will be on the surface of an accumulating screen and, with the movement in the specimen bulk by means of etching of the surface layer, the concentrations of chemical elements contained in the condensate will be smoothly decreased down to the level of their content as an admixture in the initial screen substance. Upon the construction of concentration profiles of the distribution of chemical elements over the depth of a specimen, we found, indeed, such a character of the enrichment of surface layers of accumulating screens. It turned out, however, that, in addition to the natural surface enrichment, profiles of the distribution of chemical elements show

a clearly resolved anomalous concentration maximum located at a depth of about 0.1 to 0.4 μm (see Figs. 11.25, 11.26; Sec. 11.2.6). Sometimes, we also observed the second weaker concentration maximum at a distance of 5 to 7 μm from the screen surface. It is obvious that the concentration maximum on a profile of the distribution corresponds to an enriched layer inside the bulk of the accumulating screen. We note that this layer was absent prior to the explosive experiment, because the screens were manufactured from chemically pure metals with a homogeneous distribution of admixtures.

We describe some properties of this enrichment. Most significant consists in the fact that the concentration maxima are at the same depth from the surface for all chemical elements possessing the anomalous distribution in the bulk of screens. In other words, the anomalous subsurface enrichment is not chemically stratified. Of importance also is the circumstance that its power exceeds considerably the power of the layer of a surface condensate (see the area under the corresponding segments of the curves in Fig. 11.25, Sec. 11.2.6). In an anomalous layer, we observed also chemically inhomogeneous sections. These are the aggregates of finely divided multicomponent regions which are close in composition and possess any form (see Fig. 11.27, Sec. 11.2.6). In other words, in the scope of such single aggregate, we have a spatially identical distribution of a group of chemical elements.

It is also necessary to note that the indicated regions have, most probably, the same crystal lattice as the solid solution of the matrix. Otherwise, they would acquire a crystallographic faceting or a rounded form depending on the surface energy anisotropy for various crystallographic planes of the new phase, required by the need to minimize the energy of internal interfaces upon their formation. We also add that, in order to form the new phase characterized by an increased content of a number of chemical elements as compared to the matrix, the diffusion transport of a substance across the interface from the surrounding matrix is needed. Otherwise, when the new phase is absent, the situation will be inverse: the processes of diffusion will lead only to the dissolution of the indicated regions in the surrounding matrix due to the presence of a gradient of concentrations. This yields that the origin of the aggregates under consideration cannot have the diffusion nature.

We also consider the geometry of the anomalous layer. It is lens-shaped and is located under the surface in the bulk of an accumulating screen in such a way that the screen symmetry axis passes through the lens center. That is, the depth of the layer is maximum at the center of the accumulating screen and decreases with approaching its periphery (see Figs. 11.24, 11.26; Sec. 11.2.6). This fact is not affected even by the circumstance that the central region of accumulating screens usually contains a crater: a cavity of

about 5 to 7 mm in diameter and 10 to 14 μm in depth which is formed due to the withdrawal of the screen substance as a result of its heating by the X-ray emission from the collapse zone and the ion etching by a hot plasma of the target substance dispersed after the explosion. It is remarkable that the anomalous layer repeats very exactly the form of the surface of this crater. A detailed study of the dependence of its depth on the distance from the center of the accumulating screen allows us to reveal the following regularity which is a stronger assertion: if we join the center of the zone of regeneration of the target substance with an arbitrary point of the anomalous layer by the segment of a straight line, then the length of a part of this segment which is in the bulk of the accumulating screen, l , does not depend on the distance from the point to the axis of the accumulating screen r . This yields that the depth of the layer at the distance r from the screen center $h(r)$ is connected with l by the following relation:

$$l = h(r)/\cos \vartheta = \text{const.} \quad (9.15)$$

Here, ϑ is the angle of dispersion of the target substance reckoned from the target symmetry axis coinciding with the symmetry axis of the accumulating screen. At $r = 0$, Eq. 9.15 becomes

$$l = h(0) = \text{const.} \quad (9.16)$$

That is, the constant on its right-hand side is the depth of the anomalous layer at the center of the accumulating screen.

We have given a rather detailed description of characteristics of the enriched layer discovered under the surfaces of accumulating screens in order to show the meaning of its anomaly. Discussing this question, we indicate, first of all, the circumstance that the contemporary physical science of materials does not know the mechanisms of formation of chemical inhomogeneities of a similar type. In other words, the anomalous enriched layer is not something unordinary, but it is the observation, whose origin cannot be explained, in principle, in the framework of the known physical ideas. Indeed, as the most likely mechanisms of transport of a substance across the surface in the bulk of a solid, we may consider the processes of diffusion and ion implantation (the introduction of atoms into the surface layer of a solid by means of the bombardment of its surface by ions). But it is easy to verify that any attempts to connect the origin of enriched layers with the indicated processes are faced with insurmountable difficulties.

Let us assume that the formation of anomalous layers is caused by the processes of diffusion. The penetration depth of an admixture from the surface in the solid bulk can be estimated as

$$h = \sqrt{2Dt}, \quad (9.17)$$

where D is the diffusion coefficient of atoms of an admixture in the matrix of the accumulating screen, and t is the diffusion duration.

At premelting temperatures of Cu (Cu is the screen substance, and T_{melt} is about 1083°C) for the substitution admixtures (for the majority of chemical elements, except for some elements from the first and second periods of the Periodic table), the values of diffusion coefficients are at the level of about 0.1×10^{-14} to $1 \times 10^{-14} \text{ m}^2/\text{s}$ (see Ref. 157). Eq. 9.17 yields that they need about 0.1 to 1 s and 100 to 1000 s, respectively, to penetrate at the depths of $0.2 \mu\text{m}$ and $5 \mu\text{m}$ into the screen bulk. It is seen from this estimate that the admixture cannot penetrate at the depth of the second anomalous layer at the indicated temperature by diffusion, because it is obvious that the surface of the accumulating screen has no such high temperature for this length of time. However, if we take a more realistic value for the effective temperature of the surface of the accumulating screen during the diffusion of the condensate of a plasma bunch, e.g., about 600°C , then we will get values of 100 to 1000 s and 10^5 to 10^6 s for the corresponding time intervals at the diffusion coefficients of the order of 0.1×10^{-17} to $1 \times 10^{-17} \text{ m}^2/\text{s}$ (see Ref. 152). It is obvious that we have obtained a rather overestimated value of the time interval also for the formation of the first anomalous layer with regard to the fact that Cu has a very high heat conductivity.

The insufficient time for the surface of the accumulating screen to be at high temperatures is not a single point which does not allow one to connect the formation of anomalous enriched layers with the processes of diffusion. The second point is the position of maxima on the concentration profiles of all chemical elements at the same depth from the surface of accumulating screens, i.e., the absence of a chemical stratification. The diffusion processes cannot form similar concentration profiles in principle. This is caused by the circumstance that the atoms of chemical elements differ strongly by their diffusion mobility defined by their diffusion coefficients. Indeed, at the above-mentioned premelting temperature of Cu, the values of diffusion coefficients of atoms of such chemical elements as, for example, H, B, C, N, and O in Cu exceed those of Ta, W, Pt, Au, Pb, and Bi by 5 to 6 orders (see Ref. 152). The last fact with regard to Eq. 9.17 means that the depth of the diffusion penetration of atoms of the chemical elements under study at the indicated temperature from the surface into the bulk of a Cu screen must differ by 3 orders. At a temperature of 600°C , the described situation will look more critical due to a great difference in the diffusion coefficients. In other words, concentration profiles of the diffusion nature must be chemically stratified.

Finally, we indicate one more circumstance contradicting the diffusion nature of the registered concentration profiles. The main equation describing the process of diffusion is given by the Fick law which can be presented as

$$\vec{j} = -D \operatorname{grad} C. \quad (9.18)$$

According to this law, the vector of diffusion flow density $\vec{j}$ is directed against the gradient of concentration C , i.e., the diffusion mass transfer is always realized in the direction from greater concentrations to lower ones. This means that a concentration maximum cannot be generated in principle by the processes of diffusion mass transfer. Moreover, the situation is inverse: if a concentration maximum has been formed in the bulk of a solid in some way, the diffusion processes occurring at sufficient temperatures can lead only to its “spreading” with a subsequent disappearance. In other words, the very fact of the presence of maxima on the curves of concentration profiles contradicts the assumption about their diffusion origin. For the above-mentioned reasons, upon the diffusion of an admixture from the surface into the bulk of a solid, the typical form of concentration profiles would be a monotonously decreasing curve without any local peculiarities which has the form of a “step” with a gently sloping edge expanding in the bulk of a solid. We also add that the concentration maximum includes not only the very enriched layer, but also the above-discussed chemical inhomogeneities in it, namely, the aggregates of finely dispersed regions with the enhanced content of a number of chemical elements. Hence, all the above-written relative to the enriched layer is also applicable to such aggregates.

We now consider the possibility for enriched layers appearing in the bulk of accumulating screens due to the processes of ion implantation. Some characteristics of enriched layers give evidence in the favor of just this scenario. Indeed, it is pointed out, first of all, by the symmetric position of the enriched layer relative to both the target axis and the accumulating screen and by its lens-like form. It is easy to guess that the line connecting the center of the collapse zone of a target substance with any point of the layer (see Fig. 11.24, Sec. 11.2.6) is nothing else but the trajectory of a particle of the expanding plasma bunch, and its part with length l in the bulk of the accumulating screen is the projective path (the projection of the trajectory on a direction of the initial motion of a particle) or the deceleration path of a fast particle in the bulk of the screen. In the context of the proposed interpretation of the mechanism of formation of the enriched layer, the discovered independence of l on the dispersion angle of particles ϑ or the distance from the taken point to the axis of the accumulating screen, which yields relation Eq. 9.15, looks quite logical. It is obvious that the fulfillment of this condition requires the isotropy of both the dispersion of the

target substance and properties of the substance of the accumulating screen. The first occurs by virtue of a symmetry of the explosion, and the second does due to the polycrystalline structure of the screen substance. That is, Eq. 9.15 defines naturally both the penetration depth of the front of a dispersed plasma bunch in the screen bulk and, hence, the lens-like form of the enriched layer and its position relative to the screen axis.

Moreover, the implantation mechanism of formation of the enriched layer in the bulk of the accumulating screen, as distinct from the diffusion mechanism, agrees well with the presence of maxima on concentration profiles. Changing the energy of ions in the process of ion implantation, we can get the distribution of an introduced admixture over depth to be practically of any desired form (see Refs. 134, 136). Concentration profiles with a maximum are derived upon a fixed voltage accelerating the bombarding ions, i.e., upon the typical regimes used in the implantation-based metallurgy. In this case, their energy must be sufficient in order that the layer enriched by the introduced admixture can “separate” itself from the surface and can move to the bulk of a bombarded solid. For all the ions of chemical elements, this condition is fulfilled at energies considerably exceeding 10 to 20 keV which is the upper boundary of the range of low-energy ions (see Refs. 153, 154). In the described case, the position of a maximum on the concentration profile is defined by the mean projective path l of an ion, and its width by the mean square variance of paths.

Thus, we have convinced ourselves that the implantation model of the formation of enriched layers in the bulk of accumulating screens describes quite adequately the whole series of their observed features. However, in order to accept it, we must demonstrate that it does not lead to a chemical stratification of enriched regions and describes their experimentally registered depth with the available parameters of the ion component of a plasma bunch. The analysis of spectra of the optical emission of a plasma bunch (see Chapter 4) in the explosive experiments with various targets shows that the dispersion velocity of the ions of all registered chemical elements belongs to the interval 1×10^7 to 4×10^7 cm/s. That is, they move with rather close velocities. The indicated velocity interval corresponds to the range of their kinetic energy from several keV to several tens of keV. In this case, for the ions of most chemical elements, it does not exceed 20 keV and can reach 40 to 50 keV only for the ions of two to three heavy chemical elements.

The last means that the ions of a plasma bunch are referred to low-energy ions by the criteria of implantation-based metallurgy (see Refs. 153, 154). That is, their energy is insufficient in order that the enriched layer generated by them be separated from the surface, move to the bulk of the accumulating screen, and form a maximum on the concentration profile.

Thus, the observed ions of a plasma bunch on concentration profiles are related to the surface enrichment (see Figs. 11.25 and 11.26; Sec. 11.2.6). To obtain a feel of how insufficient is the kinetic energy of the ions of a plasma bunch, we present the energies for ions Be^+ , N^+ , and Pb^+ which are necessary for them to penetrate into a Cu screen at the required depth of $0.3\text{ }\mu\text{m}$ and to form a concentration maximum there. These energies should be 0.23, 0.67, and 80 MeV, respectively (see Table 11.3, Sec. 11.2.6).

Thus, we find ourselves in a paradoxical situation: the implantation-based model is the single one which can satisfactorily describe the formation of enriched layers in the bulk of accumulating screens as a result of the explosion of targets. However, the expanding plasma bunch does not contain ions which would possess the suitable characteristics needed for the formation of the observed maxima on concentration profiles according to the requirements of this model. There is only one exit from this dead end: we should assume that the corpuscular component of a plasma bunch contains, in addition to the observed part, a part unobservable by the used spectral analytical methods. Of course, this part should be represented by superheavy chemical elements, since the chemical elements of the known part of the Periodic table are observed by the normally used methods. We note that this assumption is natural. Indeed, we may propose the other statement of the problem of searching for the exit from the situation. For example, let us try to clarify which properties should the objects contained in a plasma bunch possess, in order that they can generate, in the framework of the implantation-based model, enriched layers with the observed characteristics in the bulk of accumulating screens. Following this way, we draw easily again a conclusion that these objects are superheavy chemical elements.

First, we consider the question of why the chemical elements contained in the enriched layer cannot belong to the composition of the ion component of a plasma bunch. This assertion follows, first of all, from the fact that the ions of chemical elements of the indicated sorts in the plasma bunch had an insufficient energy to penetrate the screen bulk at a depth of several tenths of microns (see Sec. 4.2). Then we find that some other objects contained in the plasma bunch have penetrated the screen bulk to this depth and generated the whole totality of chemical elements observed as the enriched regions. The last means that, in order to form an enriched layer, nuclear transformations must have occurred. It is obvious that ordinary chemical elements do not participate in nuclear reactions under the considered conditions. Hence, the objects generating the enriched layers must be beyond the scope of the known part of the Periodic table.

In order to be in the scope of the traditional ideas of the mechanism of formation of the enriched layers, we can try to assume that the data on

the energy distribution for the ions in a plasma bunch derived in Sec. 4.2 are erroneous. For example, we assume that the energy distribution of ions in a plasma bunch was bimodal and its high-energy part was “omitted” by some reasons upon the analysis of the optical spectra of their emission. However, to penetrate the screen bulk to the prescribed depth, the ions of each sort generating the enriched layer must have the strictly defined energies in the plasma bunch which are strongly different, one from another. This requirement is conditioned by the necessity to ensure the absence of a chemical stratification in the enriched layer. The character of the energy distribution of these ions can be estimated by the data presented in Table 11.3 (Sec. 11.2.6).

Let us assume the improbable thing: the regularities of processes in the collapse zone of the target substance are such that they induce the formation of a certain sort-dependent distribution of fast ions over energy in a plasma bunch. In this case, it turns out by chance that this energy distribution of fast ions satisfies exactly the requirement for any sort of ions to penetrate to some identical depth in the copper screen. However, the character of the sort-dependent distribution of ions over energy depends very strongly on the depth, at which they generate a chemically unstratified enriched region in the copper screen. Then, based on the fact that the enriched layers in the copper screen are found at various depths (see Figs. 11.25 and 11.26; Sec. 11.2.6), we may conclude that the character of a change in the sort-dependent distribution of fast ions over energy for all the schemes of the realization of explosions should be as follows: for any of the derived distributions, the condition of the absence of a chemical stratification of the enriched region generated by it should be satisfied at a certain depth in the copper screen.

We pay attention to the circumstance that then the indicated regularity must be observed only on the Cu screens and not on any others. This follows from the fact that the dependence of the path of fast ions in a substance on the mass of atoms of its matrix is essentially nonlinear (see Refs. 153–156). That is, the equal deceleration paths of ions of different sorts in a Cu matrix become different in some other matrix of a screen. But the last assertion contradicts the experimental data: the chemically unstratified enriched regions were registered in the bulks of screens made of different materials (Al, Cu, Nb, Ta, Au, etc.). Thus, the inaccuracy of the above-presented assumptions and the contradiction following from them yield that only the particles slightly separated by sorts (by the atomic mass, rather than the nucleus charge) and with close energies would have the same path length in any matrix. We note that the formation of multicomponent enriched regions in the bulk of screens by identical particles requires again

the presence of nuclear transformations. In other words, trying to remain in the framework of the classical ideas, we have arrived again at the necessity to accept the hypothesis about the presence of superheavy chemical elements in the plasma bunch and about their key role in the formation of enriched layers.

An analogous conclusion follows also from the analysis of the other geometric peculiarities of anomalous layers. We have mentioned above that the central region of an accumulating screen contains a crater which is rather deep as compared to their depths. In this case, we have noted that the anomalous layer repeats very exactly a form of the surface of this crater and has the maximum depth in its central region. At first glance, this circumstance looks quite obvious because the implanted particles move here practically along a normal path to the surface and, hence, should have the maximum penetration depth. However, the obviousness of this fact disappears at once, if we remember that there occur not only the implantation of ions under the ion bombardment of the surface of a solid but also the intense sputtering of the very surface (see Refs. 153–156), due to which, to a great extent, the crater is formed.

Indeed, it looks reasonable at first glance that our one-type particles must be moving with a significantly higher velocity than that of the ions of the observed part of a plasma bunch. Their high velocity is, as if, necessary for the attainment of a great depth on the path in the bulk of the accumulating screen. In this case, they are the first to reach the plane screen surface and to form an enriched layer at a depth of several tenths of microns, and then the rest of the ions reach the screen surface. This induces the sputtering of the screen substance and a lowering of the screen surface. If the distance, in which the surface moves, exceeds the initial depth of the enriched layer, it will be sputtered. Otherwise, the surface only approaches the anomalous layer. Hence, at great velocities of one-type particles, the enriched layer in the central part of a screen should be absent or have a considerably lesser depth. However, this assertion contradicts the experimental observations.

It is easy to understand that the velocity of particles generating an anomalous layer cannot also be equal to that of the ions of the observed part of a plasma bunch. In this case, the processes of formation of an enriched layer and sputtering of the screen surface would have to occur simultaneously. The lowering of the screen surface induced by its ion etching will lead to that the formation of an enriched region will occur at greater depths at every successive time moment than at the previous times. This is related to the circumstance that the depth where an enriched region is formed should be reckoned, in the indicated case, from the current position of the screen surface. As a result of the running of the mentioned processes at the screen

center, we will observe both a significant increase in the width of the enriched layer and a lesser distance from the screen surface to its trailing edge. We also note that if the thickness of an etched layer will reach or exceed the penetration depth of one-type particles into the matrix, then the enriched layer will range in the screen bulk beginning from the very surface of the crater, i.e., its trailing edge will coincide with the screen surface. The assertion advanced at the beginning of this item follows from the fact that the described features of the geometry of enriched layers are not observed in experiments.

Finally, it remains to consider the last case where one-type particles generating the anomalous layer possess a velocity significantly lower than that of the ions of the registered part of a plasma bunch. In this case, the situation is very simple: first, a crater is formed due to the sputtering of the screen surface, and then the formation of an anomalous layer occurs. In this case, it is obvious that its depth should be reckoned from the final position of the screen surface. Hence, in the case under consideration, the anomalous layer should repeat exactly a form of the crater and have the maximum depth in the central region of the screen. Just these geometric peculiarities of the enriched layers were observed in explosive experiments.

Thus, we have arrived at the paradoxical conclusion: the velocity of one-type particles generating the enriched regions in the bulk of screens must be significantly lower than that of the ions of the registered part of a plasma bunch. It is obvious that these unknown one-type particles cannot be ordinary ions, since they would have a low energy by possessing a low velocity. However, at their small energy and velocity, they would fail to penetrate inside the screen to the necessary depth. Hence, they must be unusual particles, and their originality should consist, first of all, in that they must have huge energies in order to possess great paths in solids. But, at small velocities, such huge kinetic energies of particles can be ensured only at the expense of their great masses. In other words, based only on the analysis of various geometric features of the enriched layers, we infer for the third time that the unknown one-type particles generating these layers in the bulk of screens should be superheavy chemical elements.

In conclusion, we recall that the assertion about the formation of anomalous enriched regions in the bulk of accumulating screens by both the implantation of superheavy chemical elements and their nuclear transmutation is resulted from three following features of their geometry: a great depth observed, which would be impossible for the penetration of ordinary ions to achieve; the spatial coincidence of maxima on the concentration profiles (the absence of a chemical stratification); and the conjugation of a form of their surface with the crater surface upon the maximum depth at the

screen center. In this case, in order to be able to generate the observed anomalous layers, superheavy chemical elements must have the following properties: they are one-type, have low velocity, huge energy and mass, and undergo the nuclear transformations upon the interaction with atoms of a screen. These characteristics of superheavy atomic particles can be corrected by analyzing the dynamics of their deceleration in a Cu matrix (see Sec. 11.2.6). They would need to be neutral particles of a subatomic size with masses in the range 1000 to 8000 a.m.u. and have the initial velocities of about 2×10^6 to 4×10^6 cm/s. Such values of their masses and velocities define the elastic character of the interaction of their nuclei with atoms of the screen substance and, at the same time, are sufficient to ensure the relatively long paths in its bulk material.

As for a type of transformations occurring with the nuclei of superheavy chemical elements upon the formation of enriched layers, they can be referred, as noted above, to either decays leading to the creation of the whole collection of registered chemical elements, or to the stage-by-stage synthesis of heavier nuclei by means of the interaction of a primary superheavy nucleus with surrounding nuclei of a screen material. In the last case, the chemical elements registered in enriched places should be considered as by-products of the mentioned transformations of superheavy nuclei. At present, there is not, unfortunately, a sufficient amount of experimental data in order to reliably estimate the roles of each of the above-presented scenarios in the formation of anomalous layers and to choose one scenario over others. At the same time, we note that the elucidation of this question has not been considered by us as the main purpose of the present work.

9.2.9. Observation of Unidentifiable Mass-Peaks above 220 amu.

We now consider the results which were derived by the methods of mass-spectrometry and give direct evidence to the presence of superheavy chemical elements in target explosion products. These results are concerned with the discovery of the unidentifiable peaks corresponding to masses, which are not those of complexes and have values from the range of above 220 a.m.u. These peaks belong to the mass-spectra registered on the explosion products. This yields that the indicated masses are the ions of the nuclides of superheavy chemical elements. We begin from the data derived by secondary-ion mass-spectrometry (IMS 4f, CAMECA, France). The specifications of a mass-spectrometer and its operation modes realized in these investigations are described in Sec. 8.1.2 in detail. To search for the masses of superheavy nuclides, accumulating screens of the standard size with deposited target explosion products were used as specimens (see Fig. 8.1, b). The object of

the study was the deposited layer of solid products of the target explosion. In this case, to execute the sampling for the analysis, the raster of an ion microprobe (the analyzed area) was positioned directly on its surface. The procedure of a proper decoding of mass-spectra and a testing of the analyzed mass-peaks whether they belong to complex ions were carried out by the scheme which was employed upon the analysis of the isotope composition by secondary-ion mass-spectrometry and is thoroughly described in Sec. 8.1.2.

The search for the unidentifiable mass-peaks needing to be referred to the ions of the nuclides of superheavy chemical elements was performed with the use of a mass-spectrometer IMS 4f in the range of masses from 220 to 480 a.m.u. and consisted in the separation of the mass-peaks associated with complex ions from the whole set of the peaks registered in this range. The separation of the peaks of monoatomic masses in the indicated range is a rather complicated analytical problem for a number of reasons. One of the reasons is small amounts of a substance associated with the mentioned peaks, which did not allow us to carry out fully the procedure of analysis of the tested peak in all the cases before the sputtering of the studied layer of a substance. Another reason making the problem of testing of unidentifiable peaks complicated, is a considerable decrease in the resolving power of a mass-spectrometer in a range of great mass numbers. This problem and the frequent cases of the absence of any reference mass-peaks in the vicinity of the studied mass, both are deleterious to the reliability of the testing procedure.

Upon the study of the nucleosynthesis products derived in the experiments with various "target/screen" compositions, we registered most often the unidentifiable peaks corresponding to the mass numbers 271, 272, 330, 341, 343, and 433 a.m.u. We distinguish especially the peak with a mass of 433. This peak passed the analytical procedure of testing for it to be an atomic ion and revealed a special behavior which is not peculiar to the mass-peaks of ordinary monoatomic or complex ions. The last circumstance proved to be an additional weighty argument in favor of it being an ion of a superheavy nuclide. The possibility of the comprehensive study of properties of the peak with a mass of 433 was conditioned by a high yield of the secondary ion emission upon the use of accelerated Cs^+ ions as a primary ion beam. Due to this peculiarity, the peak with a mass of 433 was definitively registered on the large number of studied specimens (No. 4489, 7229, 7230, 7231, 7912, etc.).

A fragment of the mass-spectrum registered on specimen No. 7912 including the unidentifiable peak with a mass of 433 is presented in Fig. 9.27. As for the spectrum, attention should be given to its huge intensity as

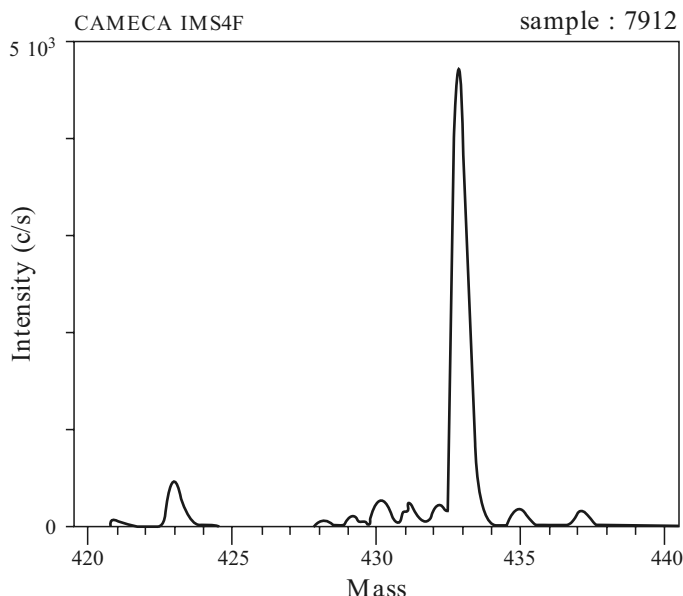


Fig. 9.27. Fragment of the mass-spectrum registered on the explosion products of a Pb target on a Cu accumulating screen and containing the unidentifiable peak with a mass of 433.

compared to those of the group of the peaks of complex ions with mass numbers 428 to 432 a.m.u. which are referred to a Pb_2O molecule. The importance of this fact is related to that the initial materials of a target and a screen in this experiment were lead and copper. The last means that the mass-peaks of their nuclides and Cs, whose ions formed the primary ion beam, would have the maximum intensities in the full mass-spectrum corresponding to the given fragment. Therefore, the mass-peaks of complexes including the nuclides of the indicated chemical elements would possess the maximum intensity. Since the peak with a mass of 433 does not obey this rule, it cannot be associated with a complex ion.

The following simple argument also indicates that the peak with a mass of 433 corresponds to an atomic ion. Indeed, while studying specimen No. 7912, we discovered that the spectrum with the 433-mass peak was not qualitatively different from several other spectra registered on close places of the surface. But the 433-mass peak was not present in the latter. The corresponding fragment of one of such spectra not containing the peak with a mass of 433 is given in Fig. 9.28. It is obvious that if the peak with a mass of 433 would correspond to a complex ion, its absence in the indicated spectra would yield the disappearance of the mass-peaks of the ions of nuclides forming this complex. Moreover, the disappearance of these peaks

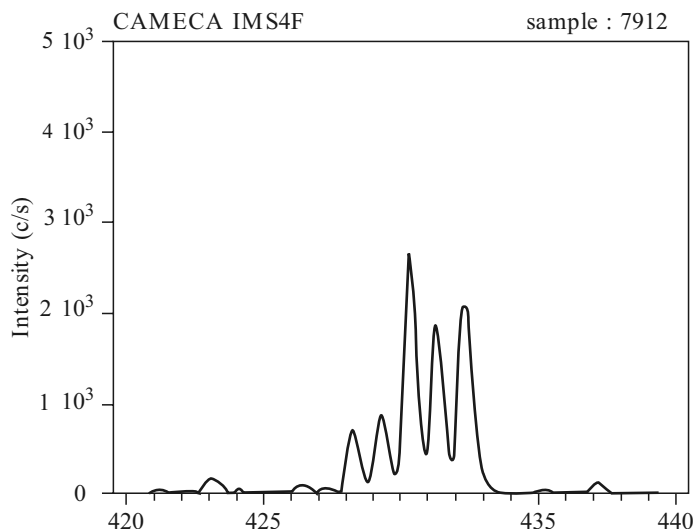


Fig. 9.28. Fragment of one of the spectra registered on specimen No. 7912 which does not contain the unidentifiable peak with a mass of 433.

will hardly be unobserved, because it would concern the most intense peaks of the analyzed mass-spectrum by virtue of the huge intensity of the peak with a mass of 433. The absence of basic differences in the collections of mass-peaks of the ions of nuclides possessing remarkable intensities in all these spectra makes this argument.

In order to clarify the nature of the peak with a mass of 433, we also used the repeated observation of the spectrum fragment containing it in the offset mode (see Sec. 8.1.2). In the majority of cases, the mentioned mode efficiently suppresses the mass-peaks of cluster ions. The testing showed that the application of the offset voltage to a studied specimen leads to a considerable decrease in the intensities of the mass-peaks of clusters of a Pb_2O molecule (see Fig. 9.27) and affects slightly the intensity of the analyzed peak with a mass of 433. The last means that the tested peak is generated by ions possessing the distribution over energies which is typical of the ions of individual nuclides, rather than those of complex ions.

The second peculiarity of the peak with a mass of 433 is its clearly pronounced local character of the distribution over the specimen surface (see Fig. 9.29). This property is very useful in clarifying the following question: Does the tested mass-peak belong to an atomic or complex ion? This question was solved by analyzing the image of the distribution of a corresponding mass on the studied area of the specimen surface. The mentioned method of identification of molecular complexes is based on the obvious fact that the image of the surface distribution of molecular mass must coincide with the



Fig. 9.29. Ion image of a fragment of the surface of specimen No. 7912 illustrating the character of the distribution of the peak with a mass of 433.

images of the surface distributions of the masses of nuclides composing it. If such coincidences are observed, then the analyzed mass-peak should be referred with a high probability to molecular ions. Otherwise, it will correspond to an atomic ion. The verification showed that the character of the surface distribution for all nuclides, being present on the studied area, does not coincide with the surface distribution of the 433-mass peak shown in Fig. 9.29. This fact testifies also in favor of the assertion that the peak with a mass of 433 belongs to atomic ions, rather than complex ones.

We would also give several other obvious arguments clarifying the nature of the tested peak. In this connection, we note that the heaviest chemical elements, for which the mass-peaks of nuclides were observed along with the 433-mass peak in the analyzed spectrum, were lead and bismuth. This means the following. In order to form a complex with the mass number equal to 433 a.m.u. from the chemical elements present on the studied area of the surface, it would be necessary to possess at least three nuclides, since two from the heaviest nuclides (^{209}Bi) cannot compose a complex with the necessary mass. In view of the circumstances that the probability of the formation of a complex under secondary emission and high vacuum drops with increase in the number of nuclides in the complex (see Ref. 99) and the tested peak reveals a huge intensity, it is reasonable to consider that the mentioned complex should contain at most three nuclides.

It is obvious that a type of this complex cannot correspond to the formula A_n , where A is the notation of a chemical element and n is a natural number or zero. The last assertion follows from the fact that 433 is a prime number. Hence, our complex should consist of the nuclides of two or three

different chemical elements, and its type should be described by one of the formulas: A_nB_m or $A_nB_mC_l$. With regard to the last formulas and the fact that there are no other unidentifiable peaks in the nearest vicinity of the peak with a mass of 433, we must inevitably accept that all chemical elements A , B , and C , which belong to the hypothetical complex, would have to be, indeed, fabricated of one-isotope chemical elements. In this connection, we note that the well-known part of the Periodical table includes only 20 stable one-isotope chemical elements. It becomes clear from the above-written that the probability of the formation of triple complexes from one-isotope chemical elements is quite low. It will be lower if we take into account that only 5 from 20 mentioned chemical elements are sufficiently heavy in order to form the necessary mass of the complex under consideration from the masses of three elements. Such chemical elements are Tb, Ho, Tm, Au, and Bi.

In this case, we may not use the notions of probability theory, but can directly verify the possibility of the formation of the considered triple complex on the basis of the indicated group of five one-isotope chemical elements. Moreover, we note that the mass-peaks of three lanthanides in the spectrum containing the peak with a mass of 433 were absent, and the mass-peaks of gold and bismuth had intensities considerably less than the intensity of the very tested peak. We also add that, as distinct from the 433-mass peak possessing a clearly pronounced local character of the surface distribution (see Fig. 9.29), the nuclides of gold and bismuth were scattered over the studied surface area more or less uniformly. Thus, the above-presented discussion and facts testify to that the 433-mass peak cannot be a complex of the nuclides of the known chemical elements. We also add that the unique isolation of the unidentifiable peak in a range of great masses is a proper sign of that it corresponds to a superheavy nuclide.

The proposed approach to the testing of unidentifiable mass-peaks can be somewhat modified. For example, if the 433-mass peak is a triple complex described by the formula A_nB_m or $A_nB_mC_l$, then the analyzed spectrum must contain the mass-peaks of double complexes corresponding to the type $A_{n-1}B_m$, A_nB_{m-1} or $A_{n-1}B_mC_l$, $A_nB_{m-1}C_l$, $A_nB_mC_{l-1}$, respectively. In this case, their amplitudes should be at least not lower than the amplitude of the peak with a mass of 433. The last assertion follows from both, the experimental facts and the probabilistic considerations. We may also specify that if the atomic mass of each of the nuclides A , B , and C does not exceed 209 a.m.u., then the mass-peaks of the mentioned double complexes should be searched for in the range of 433 to 224 a.m.u. It is obvious that the strategy of search for the mass-peaks of double complexes can be improved if we remember that they can be separated from the peak

with a mass of 433 by a mass equal to the atomic masses of the nuclides of one-isotope chemical elements. As for the application of this method to the solution of the discussed question, we note that the analysis of the considered mass-spectrum in the range from 224 to 433 a.m.u. showed the absence of any peaks with an amplitude comparable with that of the peak with a mass of 433. In other words, all the facts indicate that the 433-mass peak is a superheavy nuclide.

Finally, we consider the above-mentioned question about a particular behavior of the peak with a mass of 433 under the action of a beam of primary Cs^+ ions. For the first time, it was observed on specimen No. 7229 upon the examining of its surface distribution on the display of the monitor of an ion microprobe. We saw that the light-colored regions of the image corresponding to the aggregates of nuclides with a mass of 433 change the level of their brightness from the initially high to a slight one during the process of sputtering of the surface. It is worth noting the circumstance that if the sputtering of the surface was terminated and then was renewed after some time, the brightness of the light-colored regions was reestablished and then decreased upon further etching. This phenomenon is illustrated in Fig. 9.30 by two series of photos of the surface distribution of the peak with a mass of 433 which were made successively in some time interval during the process of sputtering of the surface. The end of the first series (Fig. 9.30, a) and the beginning of the second one (Fig. 9.30, b) are separated by a 3-min interval, during which the primary ion beam was switched off.

We note that the anomalous behavior of the nuclides with a mass of 433 under the action of the primary beam of Cs^+ ions was observed also on other specimens. This phenomenon was studied more comprehensively

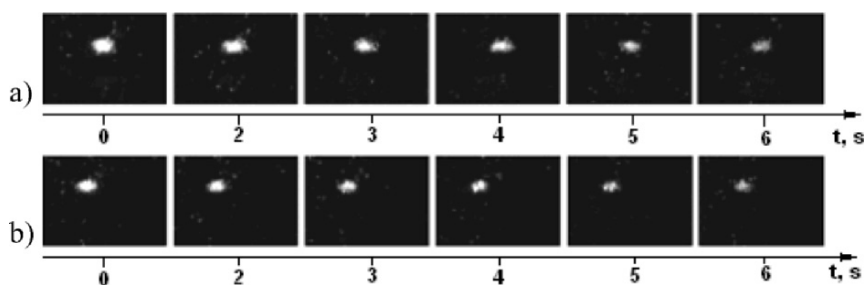


Fig. 9.30. Variation in the luminescence brightness of the regions with aggregates of the nuclides with a mass of 433 on the images of their surface distribution versus duration of the sputtering of the surface. The time interval between the end of the first cycle of etching (a) and the beginning of the second one (b) is 3 min.

on specimen No. 7912. In terms of the dependence of the secondary ion current of the studied nuclides on parameters of the primary ion beam, it consists in the following. Upon the switching-on of the primary ion beam, the process of sputtering of the surface is characterized by a certain value of the initial current of the secondary ion emission of the 433-mass nuclides which decreases with time and approaches some saturation level. The surface sputtering termination and its subsequent renewal after some time interval led to the reproduction of the described character of a variation of the secondary ion current with time. However, we observed a dependence of both the initial secondary ion current and its saturation level on the time interval, during which the primary ion beam was switched off. Moreover, the values of both quantities grew with increase in the time interval between two subsequent cycles of the sputtering of the surface. This behavior of the secondary ion current of the 433-mass nuclides is presented in Fig. 9.31.

We also found that the secondary ion current of the nuclides with a mass of 433 at the beginning of the next cycle of sputtering of the surface depends on the current of a primary ion beam on the previous cycle of sputtering. An increase in the latter led to a growth of the initial current of the secondary ion emission and did not affect its saturation level. This is

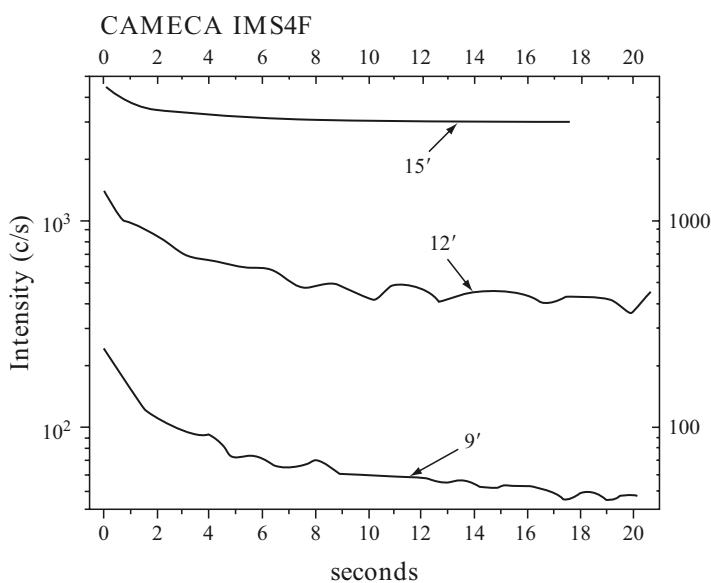


Fig. 9.31. Secondary ion current of the 433-mass nuclides versus duration of the sputtering for the time intervals between subsequent etching cycles equal to 9, 12, and 15 min.

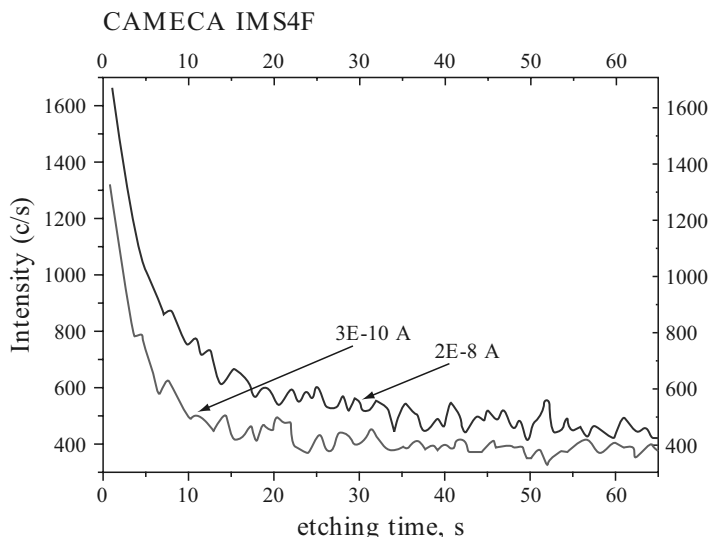


Fig. 9.32. Secondary ion current of the nuclides with a mass of 433 versus duration of the sputtering for the currents of a primary ion beam on the previous sputtering cycle equal to $3 \cdot 10^{-10}$ A and $2 \cdot 10^{-8}$ A upon the same time interval between two subsequent sputtering cycles.

illustrated by the plots of the secondary ion current of the nuclides with a mass of 433 as a function of the sputtering duration (see Fig. 9.32). In one case, the current of a primary ion beam of Cs^+ on the previous sputtering cycle was $3 \cdot 10^{-10}$ A (Fig. 9.32, lower curve). In the second case, it equaled $2 \cdot 10^{-8}$ A (Fig. 9.32, upper curve). We also mention that, in both cases, the time interval between two subsequent sputtering cycles was identical and equaled 6 min.

On the one hand, we consider the above-described behavior of the nuclides with a mass of 433 under the action of a primary beam of Cs^+ ions as anomalous, because it cannot be explained by the influence of the changes in properties of the matrix conditioned by the implantation of ions of the primary beam and by the adsorption of active admixtures by the surface on their secondary ion current, since the manifestation of the mentioned effects has other character (see Refs. 102–104). On the other hand, the indicated behavior was not characteristic of chemical elements of the known part of the Periodical table, which were present on the studied areas of the surface of specimens, and is inherent only in superheavy nuclides associated with unidentifiable peaks in a range of great masses.

The registered character of the dependences of the secondary ion current of nuclides possessing a mass of 433 a.m.u. On the conditions of sputtering is such that it can be explained, in our opinion, only by

the occurrence of nuclear reactions in the substance of target explosion products under the action of the beam of an ion microprobe. The possibility of the running of such processes under the indicated conditions is testified by the results of experiments on the registration of the decays of superheavy chemical elements which were initiated by low-energy perturbations (see Chapter 10). Moreover, the presence of nuclear transformations in the case under study is uniquely indicated by the discovered effect of accumulation of the nuclides with a mass of 433 in the intervals between the subsequent sputtering cycles of the surface (see Figs. 9.30–9.32). By virtue of the fact that the observed spectra do not contain other unidentifiable mass-peaks, it is reasonable to consider that the accumulation of these nuclides occurs due to the decay of heavier nuclei, whose masses are beyond the range of analyzed masses of a mass-spectrometer IMS 4f.

While analyzing the kinetics of the indicated transformations, it is convenient to use the data on the secondary ion saturation currents. As distinct from initial secondary ion currents, they are defined more exactly, on the one hand, but depend, on the other hand, actually only on one parameter: the time interval between two subsequent sputtering cycles (see Figs. 9.31 and 9.32). As for the fact of the excess of initial secondary ion currents over the corresponding secondary ion saturation currents, it is obviously conditioned by a larger degree of the perturbing action of the primary ion beam on superheavy nuclides subjected to a stimulated decay which are located closer to the surface of a specimen under study.

For the sake of generality, we will consider the case where a radioactive decay leads to the formation of atoms which are, in turn, also radioactive, i.e., where a chain of decay is realized. We assume that an atom of substance B is transformed into an atom of the substance A , and an atom of substance B appeared, in turn, due to the decay of an atom of substance C , etc. We may extend this chain in both sides up to any number of links. In the scope of a link of the chain, the decayed and formed atoms relative to each other are, respectively, parent and daughter ones. Decays of such a kind are called successive ones, and the groups of nuclides participating in them are called radioactive families. The analysis of the kinetics of a successive decay is actually reduced to solving the problem of determination of the temporal dependence of the amounts of substances A , B , C , etc. We denote these amounts by the symbols of the corresponding substances. Then, to describe their dependence on time, we will use such parameters as the mean lifetime of a radioactive nuclide (τ_A , τ_B , τ_C , etc.) and the decay probability (λ_A , λ_B , λ_C , etc.) (see Refs. 157, 158). The physical sense of these parameters yields the obvious connection between them: the time τ is equal to the reciprocal of a decay probability λ .

Let A be the substance of a 433-mass nuclide, and let B be the substance of its parent nuclide. Then the change of the amount of a parent substance is described by the relation

$$\frac{dB}{dt} = -\lambda_B B. \quad (9.19)$$

Per unit time, $\lambda_A A$ atoms of substance A decay. However, at the same time, $\lambda_B B$ decays of substance B occur. Since each decay of an atom of substance B results in the creation of an atom of substance A , $\lambda_B B$ atoms of substance A is formed per unit time. Therefore, a change in the amount of a daughter substance can be presented as

$$\frac{dA}{dt} = -\lambda_A A + \lambda_B B. \quad (9.20)$$

Equations 9.19 and 9.20 form a system of differential equations and have such a view that they can be solved successively, beginning from the first one. Each time, we deal only with one equation and one unknown. This remains true also in the case of a longer decay chain or a greater number of equations in the system.

Solving the derived system of equations results in the relations

$$B(t) = B_0 \exp(-\lambda_B t), \quad (9.21)$$

$$A(t) = A_0 \exp(-\lambda_A t) + \frac{\lambda_B B_0}{\lambda_A - \lambda_B} [\exp(-\lambda_B t) - \exp(-\lambda_A t)], \quad (9.22)$$

where A_0 and B_0 are the values of $A(t)$ and $B(t)$ at $t = 0$, i.e., the initial amounts of atoms of the daughter and parent substances. In the general case, Eq. 9.22 describes a change of the amount of substance A with time. Its specific form depends on the initial conditions defined by values of A_0 and B_0 , the ratio of the parameters λ_A and λ_B characterizing the decay rate of substances A and B , and the time interval, on which the equation is considered. Let us try to select the mentioned parameters so that Eq. 9.22 will describe a change of the secondary ion saturation current of a 433-mass nuclide in the course of time (see Fig. 9.31).

The data presented in Fig. 9.31 yield that the secondary saturation current of substance A (a 433-mass nuclide) increases sharply during the decay time of substance B . For example, after 9, 21 (9+12), and 36 (9+12+15) min of the decay of a parent substance, its amounts correspond to the counting rates $6 \cdot 10^1$, $6 \cdot 10^2$, and $2.5 \cdot 10^3$ pulse/s, respectively. This means that, at the beginning of a decay, the amount of the daughter substance was small. That is, without any loss in strictness, we can set $A_0 = 0$. Then Eq. 9.22 takes the form

$$A(t) = \frac{\lambda_B B_0}{\lambda_A - \lambda_B} [\exp(-\lambda_B t) - \exp(-\lambda_A t)]. \quad (9.23)$$

Based on the same plots, we can conclude that the parent substance decays considerably faster than the daughter one. By the way, the daughter substance A is long-lived or stable in the general case. This is evidenced by the fact that substance A is registered at the initial stage of the measurement on specimens stored several months after explosive experiments. It is obvious that we deal in this case with nuclide A generated upon the explosion of a target, rather than upon the induced decay of other superheavy nuclides. That is, substance A is of the “relict” origin.

In other words, we have $\lambda_A \ll \lambda_B$ or $\tau_A \gg \tau_B$. In this case, the whole process described by Eq. 9.23 splits into two stages. First, at t being of the order of τ_B , the transformation of B into A occurs. In this case, substance A does not decay practically, because the inequality $\tau_A \gg \tau_B$ holds. Then, on the second stage at $t \gg \tau_B$, we observe the slow decay of substance A . Let us elucidate the form of Eq. 9.23 on the first decay stage. First of all, the fractional factor in Eq. 9.23 takes the form

$$\frac{\lambda_B B_0}{\lambda_A - \lambda_B} \approx \frac{\lambda_B B_0}{-\lambda_B} = -B_0 \quad (9.24)$$

for $\lambda_A \ll \lambda_B$.

We now transform the second exponent in the square brackets in Eq. 9.23. Because $t \ll \tau_A = 1/\lambda_A$, we get $\lambda_A t \ll 1$. This yields that $\exp(-\lambda_A t) \approx 1$. Thus, with regard to the above-presented comments, Eq. 9.23 at t being of the order of τ_B takes the form

$$A(t) \approx B_0[1 - \exp(-\lambda_B t)]. \quad (9.25)$$

On the second stage at $t \gg \tau_B = 1/\lambda_B$, we get $\lambda_B t \gg 1$. Here, we can neglect the small value, $\exp(-\lambda_B t)$, because it is small not only relative to 1, but relative to $\exp(-\lambda_A t)$ in view of the fact that $\lambda_A \ll \lambda_B$. In other words, on the second stage, Eq. 9.23 is transformed to the form

$$A(t) \approx B_0 \exp(-\lambda_A t). \quad (9.26)$$

In Fig. 9.33, we show the curves corresponding to Eqs. 9.25 and 9.26 by dotted lines. We note that, in Fig. 9.33, the exact Eq. 9.23 describing a change of the amount of substance A with time corresponds to a continuous smooth curve derived by sewing the initial section of curve 1 and the final section of curve 2. This sewing occurs, of course, in a time interval of the order of τ_B .

Now we may consider the question as to how adequately the smooth curve (or its separate segments) plotted in Fig. 9.33 describes the experimentally observed change of the amount of a 433-mass nuclide in the course

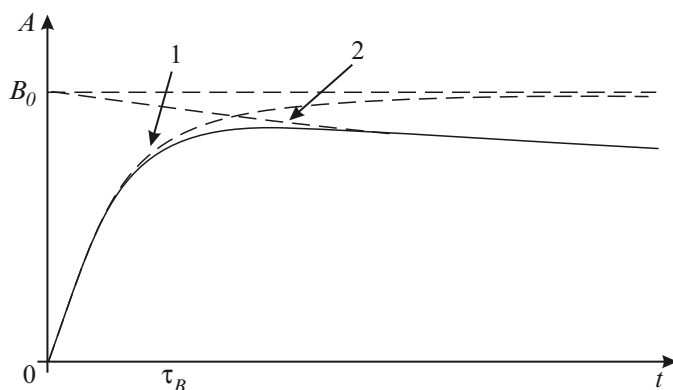


Fig. 9.33. The amount of the long-lived daughter substance A versus time in the case of short-lived parent substance B . Curve 1 describes the growth of the amount of substance A at the expense of the fast decay of substance B . Curve 2 shows a decrease of the amount of substance A at the expense of the own decay.

of time. It is obvious that its segment on the second stage does not suit us, because the amount of the daughter substance decreases there, whereas we observed experimentally a growth of the amount of a 433-mass nuclide. Let us move to the segment of the smooth curve on the first stage which shows the increase in the amount of the daughter substance A . However, this growth differs by its character from that revealed by a 433-mass nuclide. Indeed, the generation rate for 433-mass nuclides sharply increases with time (see Fig. 9.31), whereas the generation rate for nuclides of the daughter substance A decreases (see Eq. 9.25 and the first segment of the curve in Fig. 9.33). This difference is essential. By virtue of this fact, no segment of the curve in Fig. 9.33 describes adequately the temporal change of the amount of a 433-mass nuclide defined by the running of nuclear processes.

At once, we discount the possibility of the running of nuclear transformations in this case by the scenario related to the entering of nuclides B into the stage-by-stage synthesis of heavier nuclei by means of their interaction with surrounding nuclei of a screen substance as a reason for this failure. Indeed, it would be extremely striking to consider, in this case, that a by-product of such a synthesis is exclusively one nuclide, namely, a 433-mass nuclide. On the other hand, having absorbed 6-7 atoms of Cu and ejected a 433-mass nuclide as ashes, nuclide B would return in the initial state, and no stage-by-stage synthesis will occur (for the synthesis, it is necessary to absorb heavy nuclei and to emit light ones). Simply saying, Cu atoms would be reprocessed into 433-mass nuclides. With regard to the circumstances that this process occurs upon the switched off primary ion beam (i.e., it

does not need any stimulation) and its rate increases, we are led to conclude that all the substance of the accumulating screens should be reprocessed, which is not observed in experiments.

In order to understand the real reasons for the failure of this scenerio, it is necessary, first of all, to analyze the correctness of the choice of initial conditions for the system of differential equations; Eqs. 9.19–9.20. We have accepted that the parent nuclide B is short-lived and the daughter nuclide A is long-lived. This assertion arises no doubts. However, then we have silently assigned a nonzero value B_0 to the initial amount of a nuclide of substance B . We note that there were no grounds to make such an assignment. Let us analyze the meaning of the mentioned condition. First of all, if nuclide B is short-lived, it cannot be that nuclide which underwent a stimulated decay. Indeed, otherwise it would decay during the long-term storage of the specimen after the explosive experiment by virtue of its short lifetime because the storage time exceeds the mean lifetime of nuclides of the parent substance τ_B , which is about several minutes, by many orders of magnitude. Hence, in this case we would have nothing to undergo a stimulated decay in our experiments. Thus, the condition $B(0) = B_0$ assumes that we observed the stimulated decay of some superheavy nuclide C which generated the amount of nuclide B equal to B_0 during the action of the primary ion beam for several tens of seconds (see Fig. 9.31). Thus, nuclide B served as the initial parent substance for the decay under consideration. However, the scenario of the development of events can be different. For example, let us assume that nuclide D undergoes the stimulated decay which leads to the creation of the short-lived nuclide C in the amount of C_0 . In this case, there was no short-lived nuclide B in the specimen, and it appears only as a result of the decay of nuclide C . In this case, the initial amount of nuclide B is zero, but it will increase rapidly with time. This scenario is prompted by the form of Eq. 9.25. Indeed, if it would contain a rapidly growing function instead of the constant B_0 , we could describe, possibly, the observed growth of the 433-mass nuclide.

Thus, we will correct the scenario of a successive decay with regard to the above-made comments. To this end, we must find, first of all, the dependence of the amount of substance B on time. This can be done by considering the link of the successive decay corresponding to the transformation of substance C in B which is described by the system of differential equations analogous to system Eqs. 9.19–9.20:

$$\frac{dC}{dt} = -\lambda_C C, \quad (9.27)$$

$$\frac{dB}{dt} = -\lambda_B B + \lambda_C C. \quad (9.28)$$

According to the above discussion, we set $C(0) = C_0$ and $B(0) = 0$ as its initial conditions. Then its solution reads

$$C(t) = C_0 \exp(-\lambda_C t), \quad (9.29)$$

$$B(t) = \frac{\lambda_C C_0}{\lambda_B - \lambda_C} [\exp(-\lambda_C t) - \exp(-\lambda_B t)]. \quad (9.30)$$

Now, to search for a specific form of $B(t)$, it is necessary to define the ratio of the parameters λ_C and λ_B . If we set the parent substance C to be short-lived relative to the daughter substance B , we arrive at the following scenario of the successive decay. During the time interval of the order of τ_C , being much less than τ_B , practically all substance C has transformed into substance B , and only then substance B begins to transform into substance A . This actually means that $B(0) \approx C_0 \neq 0$, i.e., we meet the variant which was considered above and did not lead us to the desired result.

Thus, we turn to the case of a long-lived parent substance C and a short-lived daughter substance B :

$$\tau_C \gg \tau_B, \quad \lambda_C \ll \lambda_B. \quad (9.31)$$

Inequalities Eq. 9.31 yield, first of all, that we can neglect λ_C as compared to λ_B in the denominator of the factor in Eq. 9.30. Hence, Eq. 9.30 takes the form

$$B(t) \approx (\lambda_C C_0 / \lambda_B) [\exp(-\lambda_C t) - \exp(-\lambda_B t)]. \quad (9.32)$$

In order to understand the behavior of the quantity $B(t)$ described by Eq. 9.32, we consider two stages of the process for small and large values of t as compared to τ_B . First, we turn to the stage when the time t passed after the beginning of the process exceeds significantly τ_B . That is, $\lambda_B t \gg 1$. The last condition means that we can neglect $\exp(-\lambda_B t)$ in the square brackets. Hence, at $t \gg \tau_B$, Eq. 9.32 becomes

$$B(t) \approx (\lambda_C C_0 / \lambda_B) \exp(-\lambda_C t). \quad (9.33)$$

We should like to make some comments on the physical content of the derived solution. First of all, we note that substance B , being formed at the beginning of the process, has completely decayed by the time t under consideration. In this case by virtue of the fact that B is decaying rapidly, we have only substance B , being formed recently, at every time moment. In other words, we are faced with a stationary state in the case under study: substance B is being formed from C and is decaying immediately. In this case, substance B is not accumulated, since it decays rapidly. But substance

B does not disappear by virtue of the fact that it is constantly formed from substance C . The stationary state is characterized by the fact that the amount of substance B is practically constant, because, during a time unit, the number of decayed atoms of B is approximately equal to that of atoms of B formed from C . The above follows from Eq. 9.33. Indeed, with regard to the fact that $C(t) = C_0 \exp(-\lambda_C t)$, Eq. 9.33 can be rewritten as

$$\lambda_B B(t) \approx \lambda_C C(t). \quad (9.34)$$

However, Eq. 9.34 means just the coincidence of the numbers of decayed (the left-hand side of the relation) and formed (the right-hand side) atoms of B at every time moment. Finally, Eq. 9.34 can be written in the form

$$B(t) \approx (\lambda_C / \lambda_B) C(t). \quad (9.35)$$

With regard to the condition $\lambda_C \ll \lambda_B$, relation Eq. 9.35 yields that the instantaneous amount of substance B in the stationary state is proportional to the amount of substance C , the latter being always much more than the former.

We now consider the initial stage of the process where $t < \tau_B = 1/\lambda_B$. In this case, $\lambda_B t < 1$, and $\lambda_C t$ is still smaller by virtue of the inequality $\lambda_C \ll \lambda_B$. Hence, by expanding $\exp(-\lambda_C t)$ and $\exp(-\lambda_B t)$ in Eq. 9.30 in series and taking only two first terms of the expansion, we get

$$B(t) \approx \frac{C_0 \lambda_C}{\lambda_B - \lambda_C} (1 - \lambda_C t - 1 + \lambda_B t) = C_0 \lambda_C t. \quad (9.36)$$

It follows from Eq. 9.36 that, at the initial time moment, we have

$$B(0) \approx 0, \quad (9.37)$$

and then the amount of substance B grows linearly with time. This is just what we try to get. But the last means that, at the initial time moment, we deal with a nonstationary state. Indeed, for a stationary state (see Eq. 9.35), we would initially get

$$B_{st} \approx (\lambda_C / \lambda_B) C_0. \quad (9.38)$$

Let us estimate the time t_{st} which is required to reach the amount B_{st} at a constant initial rate of increase in the amount of substance B , i.e., to approach the stationary stage. This time can be determined from the condition

$$B(t_{st}) \approx B_{st}. \quad (9.39)$$

In view of Eqs. 9.36 and 9.38, condition Eq. 9.39 yields the following estimate for the time interval needed to approach the stationary stage:

$$t_{st} \approx 1/\lambda_B = \tau_B. \quad (9.40)$$

Thus, the stationary state is reached for the time interval approximately equal to the mean decay time of substance B . In this case, the condition $\tau_C \gg \tau_B$ yields that the amount of substance C is practically constant for this time interval.

The curves corresponding to Eqs. 9.33 and 9.36 are presented in Fig. 9.34 by dashed lines. Curve 1 describes a linear increase in the amount of substance B at the initial stage of the process, and curve 2 corresponds to the stationary stage. In this case, the exact Eq. 9.30 describing a variation in the amount of substance B with time corresponds to the continuous smooth curve in Fig. 9.34 which is drawn by sewing the initial segment of curve 1 and the final segment of curve 2 on the time interval of the order of τ_B . In Fig. 9.34, the time moment t_{st} corresponds to the abscissa of the point of crossing of the inclined straight line (curve 1) with the horizontal straight line corresponding to the level B_{st} .

We are ready now to return again to the question about the transformation of substance B in substance A . However, due to Eq. 9.36, the system of differential equations Eqs. 9.19–9.20, with which we have described the process under consideration, becomes

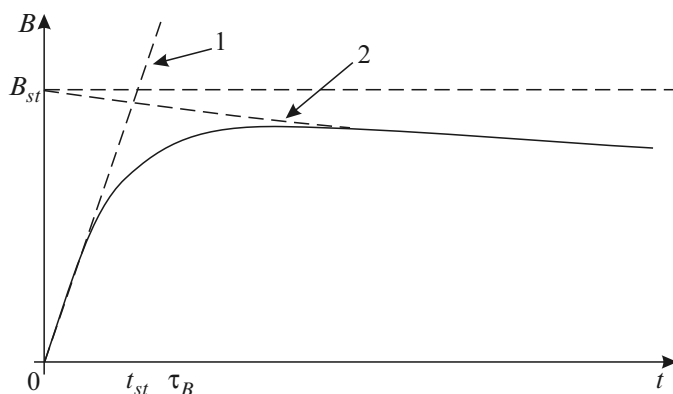


Fig. 9.34. Time dependence of the amount of a short-lived daughter substance B in the case of a long-lived parent substance C . Curves 1 and 2 describe, respectively, a linear growth of the amount of substance B at the initial stage and a slow decrease in the amount of substance B at the stationary stage.

$$B(t) \approx C_0 \lambda_C t, \quad (9.41)$$

$$\frac{dA}{dt} = -\lambda_A A + \lambda_B B. \quad (9.42)$$

It is worth noting that the new system of differential equations can be solved only in the time interval $t < \tau_B$, since Eq. 9.41 is valid only in this case. We also recall that, in turn, $\tau_B \ll \tau_A = 1/\lambda_A$. The last means that the first term on the right-hand side of Eq. 9.42 is negligible as compared to the second term. That is, Eq. 9.42 can be presented with regard to Eq. 9.41 and the above-written formulas as

$$\frac{dA}{dt} \approx \lambda_B \lambda_C C_0 t. \quad (9.43)$$

A solution of this differential equation is the function

$$A(t) \approx \frac{1}{2} \lambda_B \lambda_C C_0 t^2. \quad (9.44)$$

Thus, Eq. 9.44 describes a change in the amount of substance A with time, and Eq. 9.43 gives its growth rate as a function of time. It is easy to see that Eq. 9.44 describes the experimental growth of the amount of the daughter substance A , and Eq. 9.43 indicates that the increase in the amount of substance A occurs with an increasing rate. In other words, we get a solution which is in a good qualitative agreement with the experimental temporal behavior of the amount of the 433-mass nuclide.

We now consider the quantitative aspect of the comparison of the derived solution describing a change in the amount of the daughter substance A during the successive decay with the experimental temporal dependence of the amount of the 433-mass nuclide. This comparison can be carried out by using the obvious assertion that the secondary ion saturation current of nuclide A should be proportional to the amount of its substance, i.e., the relation

$$i_A(t) = kA(t) \quad (9.45)$$

should be valid, where k is a constant coefficient. However, in view of Eq. 9.44, relation Eq. 9.45 can be presented as

$$i_A(t) \approx \frac{1}{2} k \lambda_B \lambda_C C_0 t^2 = k^* t^2. \quad (9.46)$$

We note that the left-hand side of Eq. 9.46 includes the experimental values of the secondary ion current of the 433-mass nuclide, and its right-hand side is the analytic formula which is derived by us and describes its

change with time. It is obvious that a quantitative measure of the adequacy of the description with the proposed analytic formula for the temporal dependence of the secondary ion current could be the constancy of the coefficient k^* on the right-hand side of Eq. 9.46.

Prior to the execution of such a verification, we give explicitly the formula for k^* . Eq. 9.46 yields

$$k^* = \frac{1}{2}k\lambda_B\lambda_C C_0. \quad (9.47)$$

We also add that the coefficient k^* has a dimension of pulse/(s · min²), because the secondary ion current upon the registration of the 433-mass nuclide is characterized by the counting rate of pulses, and the time interval between two subsequent sputtering cycles, which corresponds to the free decay, is measured in minutes (see Fig. 9.31). Now we give the values of the coefficient k^* corresponding to three experimental values of the secondary ion saturation current of the 433-mass nuclide. They are about 0.7, 1.4, and 1.9 pulse/(s · min²) for the free-decay time intervals of 9, 21, and 36 min, respectively.

Analyzing the presented values of the coefficient k^* , we note, on the one hand, that they differ not very strongly. So, we may conclude that the expression on the right-hand side of Eq. 9.46 describes quite satisfactorily the temporal dependence of the secondary ion saturation current of the 433-mass nuclide. However, on the other hand, we cannot but note that values of the coefficient k^* have a clear tendency to slowly grow with time. This tendency could be ignored, if it were not for the obviousness of its origin. We should like to comment on the nature of the mentioned dependence. It becomes understandable, if we remember that the primary ion beam was switched on each time for several tens of seconds upon the registration of the secondary ion saturation current of the 433-mass nuclide. According to the above-described scenario of the successive decay, this circumstance led to the stimulated decay of the superheavy nuclide D with the formation of an additional amount of the daughter nuclide C . According to Eq. 9.47, these processes gave an additional contribution, in turn, to a value of the coefficient k^* and manifested in some increase in the registered rate of generation of the 433-mass nuclide.

The above-written discussion is clarified by the plots presented in Fig. 9.35. If the secondary ion saturation current of the 433-mass nuclide were not registered at the 9th minute, its growth would be described by a quadratic parabola with the coefficient $k_1^* = 0.7$ during the whole time interval. However, we carried out the measurement at the 9th minute, and this led to the increase in the coefficient k^* up to 1.4 (see the branch of the parabola with the coefficient k_2^* in Fig. 9.35). Therefore, during the

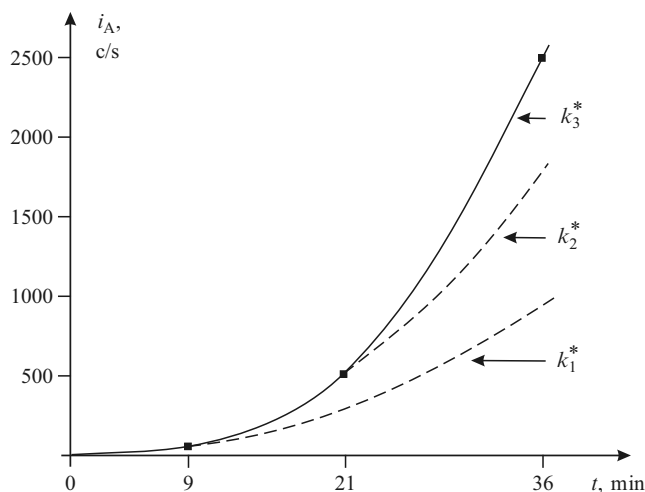


Fig. 9.35. Scheme of the approximation of the secondary ion current of the 433-mass nuclide with a piecewise quadratic function of time.

subsequent 12 min (from the 9th till the 21st minute), the growth of the amount of the 433-mass nuclide occurred by the quadratic law with $k_2^* = 1.4$. Finally, the switching-on of the primary ion beam at the 21st min increased again the value of k^* up to 1.9 (see the branch of the parabola with the coefficient k_3^* in Fig. 9.35). This yielded that the secondary ion saturation current followed the segment of a quadratic parabola with the coefficient $k_3^* = 1.9$ during the last 15 min (from the 21st till the 36th minute).

Thus, all the peculiarities of a change of the secondary ion current corresponding to the 433-mass nuclide as functions of the parameters of the action of a beam of primary ions Cs^+ , which seem firstly to be anomalous, are described satisfactorily and completely by the above-proposed model of a successive decay of the family of superheavy nuclides, whose product is the 433-mass nuclide. In the absence of alternative versions of the solution of the problem under study, the successful application of this model gives evidence, on the one hand, to the correctness of the conclusion that the peak with a mass of 433 a.m.u. corresponds to a superheavy nuclide and indicates, on the other hand, that the studied target explosion products must contain, together with the 433-mass nuclide, at least three superheavy nuclides, whose mass numbers are beyond the limits of the working range of masses analyzed by a secondary-ion mass-spectrometer IMS 4f.

Thus, by the example of one of the whole series of mass-peaks registered on target explosion products in the mass range above 220 a.m.u., we have illustrated the working procedure as for the testing whether they belong to atomic ions. This procedure consists in the separation of

mass-peaks referred to complex ions from the whole set of peaks registered in this range. By virtue of the variety of properties, analytic specificity, and some peculiarities of the behavior of the tested mass-peaks, we slightly modified the procedure sometimes, but its basic principles remain unchangeable in all the cases. The mass-peaks which passed the procedure were referred by us to the kind of basically unidentifiable mass-peaks in the scope of the known part of the Periodic table of chemical elements. The last means that all these peaks should be associated with ions of the nuclides of superheavy chemical elements, being present in the laboratory nucleosynthesis products.

9.3. Study of the Composition of Target Explosion Products by Independent Laboratories

The target explosion products were studied by various methods of mass-spectrometry also at other scientific institutions of Ukraine and at several foreign scientific-research laboratories. The subject of these studies was the determination of the isotope and element compositions of the nucleosynthesis products. In this case, a special attention was paid to the analysis of the mass-spectra registered on target explosion products in a range of high masses. In these studies, several tens of mass-peaks from the mass range above 220 a.m.u. which cannot be interpreted as those of complex ions were observed. We note that some peaks coincide with those referred by researchers of the Electrodynamics Laboratory “Proton-21” to the kind of unidentifiable ones or to the peaks of the nuclides of superheavy chemical elements. Below, we briefly comment on some of the mentioned results.

9.3.1. Comments to the Official Conclusion of the Concern “Luch”, Russia, Regarding the Objects given by our Laboratory for their Study with a Mass-Spectrometer “Finnigan” Mat-262

A series of specimens including more than 20 accumulating screens with deposited target explosion products was studied by thermoionization mass-spectrometry (MAT-262, Finnigan) at a laboratory of the Scientific-Research Institute “Luch” (Ministry of Atomic Engineering, Russia). As the specimens studied directly with a device, we took small bands of $1.5 \times 20 \times 0.1$ mm in size cut from accumulating screens. The mass-spectrometer contained a rhenium ionizator. The range of analyzed masses was from 6 to 460 a.m.u.

According to the data of measurements performed at that laboratory, a large number of unusual noninterpretable mass-peaks was observed on the studied specimens in the range of heavy masses. The inability to identify them as those of complex ions was conditioned by the following

circumstances. In some cases, the masses of the mentioned peaks did not correspond to any known combination of isotopes from the standard catalogs of masses. But even if a catalog contained the complexes with suitable mass numbers, the spectra under study did not include the necessary collection of the mass-peaks of nuclides entering the composition of the tested complex. In other words, the discovered mass-peaks had all the attributes typical of the mass-peaks of atomic ions. However, their masses exceeded those of all known ions of nuclides.

The noninterpretable peaks were grouped mainly in two mass ranges: from 253 to 292 a.m.u. and, to a lesser extent, from 350 to 440 a.m.u. Most frequently registered, were the peaks with mass numbers of 271, 272, 277, 280, 330, 341, 343, and 394 a.m.u. As an example, we present one of these peaks in Fig. 9.36. We also note that the noninterpretable peaks included the mass-peak with a mass of 433 a.m.u. which was registered many times as a unidentifiable peak by researchers of the Electrodynamics Laboratory "Proton-21" on a mass-spectrometer IMS 4f.

Upon the execution of the analysis and identification of registered masses, we used the catalogs containing the maximally complete data on molecules and clusters, whose masses correspond to the working range of the device. Among the reference books, we mention the ICP-MS Interferenze Tabelle of the Finnigan MAT firm, Lawrence Berkeley National Laboratory Mass-Reference Handbook, Internet Mass-Catalog of Moscow State University, and other Russian mass-catalogs.

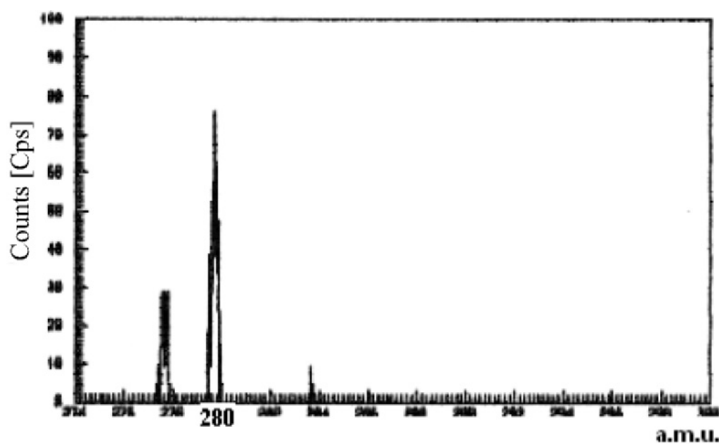


Fig. 9.36. Fragment of a mass-spectrum containing the noninterpretable mass-peak of a mass of 280 a.m.u. (accumulating screen No. 9105, MAT-262, Finnigan).

9.3.2. Comments to the Official Conclusion of United Metals LLC, USA, Report Sims-030623

Several specimens cut from the accumulating screens with deposited target explosion products were studied at the laboratory UNITED METALS LLC (USA) with the help of a mass-spectrometer SIMS 030623, where the source of primary ions was a liquid-metal gallium gun with an energy of 6 to 9 keV. The range of masses analyzed by a mass-spectrometer was 1 to 300 a.m.u. The studied specimens were fragments of accumulating screens Nos. 7753 and 7754 stored for a time more than one month after the explosive experiments.

According to the report given by the laboratory UNITED METALS LLC, the groups of peaks of the ions of heavy nuclides in the mass range above 220 a.m.u. were registered on both specimens and were not reliably identified, because they were absent in the available databases. The mentioned peaks were present in the mass-spectra of positive and negative secondary ions. The mass-peaks discovered on both specimens have the following mass numbers: 221, 222, 223, 224, 224.2, 225, 232.4, 232.8, 233.4, 235, 236, 238, 239, 240, 241, 246, 246.6, 247.2, 247.6, and 248.4 a.m.u. Moreover, the spectra which were registered on specimen No. 7753 included also the peaks with mass numbers of 243, 257, 257.6, 259, 263.2, 265.2, 266, 268, 276.6, and 284.6 a.m.u. Fragments of the spectra containing some of the enumerated unidentifiable mass-peaks are given in Fig. 9.37.

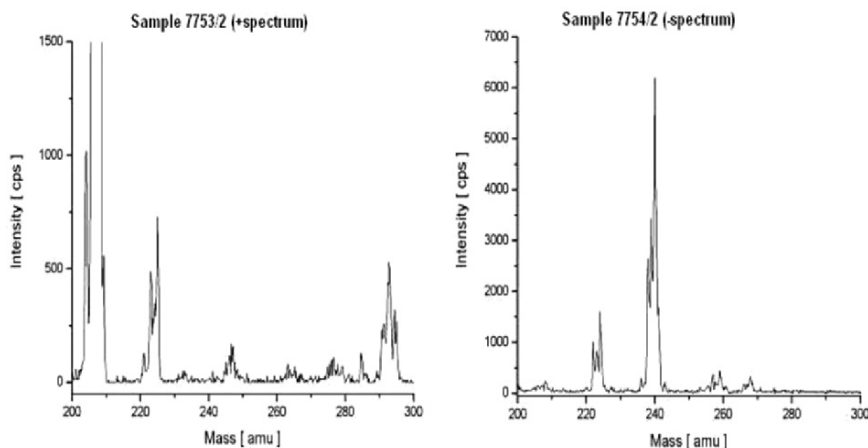


Fig. 9.37. Fragments of the mass-spectra which are registered on accumulating screens No. 7753 (a) and No. 7754 (b) and contain the unidentifiable peaks in the mass range of 220 to 300 a.m.u. (UNITED METALS LLC, SIMS 030623).

We give also the comment of researchers of the laboratory UNITED METALS LLC, which is contained in their report and concerns the identification of the observed peaks. By the mentioned reasons, they have no possibility of unambiguously associating a specific nuclide with each registered peak by indicating the mass and charge of nucleus. Nevertheless, the first group of the given mass-peaks is related to the presence of the following groups of chemical elements present in the products of explosions: Rn, Fr, Ra, Th, Pa, U, Np, Cm, and Bk. As for the second group of mass-peaks, they belong, in their opinion, to the nuclides of Fm, Md, No, Lr, Db, Sg, Bh, and Mt.

In the last comment, several points deserve special attention. First of all, striking is the fantastic chemistry of the studied specimens. It is difficult to imagine a specimen, whose composition could contain the collection of such exotic chemical elements. However, the point is not only that they are exotic in nature but also that there is an absence of modern technical abilities for producing the specimens with such a composition. Indeed, let us consider the second group of chemical elements which without any exception, have no stable isotopes. Moreover, the mean lifetime of the most long-lived known nuclides of Fm and Md is counted in days, of No and Lr – in minutes, and of Db, Sg, Bh, and Mt – in seconds and fractions of seconds (see Refs. 159, 160). With regard to the last circumstance and the fact that the specimens were stored at least one month prior to the investigation, it is easy to conclude that at least No, Lr, Db, Sg, Bh, and Mt would completely decay and disappear from the specimens by the time moment of their observation. It is possible to propose two versions in order to explain their presence in the specimens: these chemical elements have also other nuclides which are unknown up to now and are stable but for a fast decay, or they appear during the observation due to the decay of heavier nuclides contained in the specimens which is stimulated by the primary ion beam.

9.4. Main Results and Conclusions

Let us summarize the discussion on the presence of superheavy chemical elements in the laboratory nucleosynthesis products. We have presented a sufficiently large number of direct and indirect experimental facts testifying to their production in the explosive experiments. These facts were derived as a result of the diagnostics of both the collapse zone of a target substance and a dispersed plasma bunch and upon the study of the composition of products of the explosions, which were deposited on accumulating screens, by various analytic methods. We recall that these facts, first of all, include: unidentifiable X-ray and Auger-peaks; centralized track clusters with an anisotropic distribution of tracks; initiation of the emission of nuclear particles by

low-energy perturbations; violation of the energy balance, which is indicated by the calculations of the energy yield of running nuclear reactions by using the results of the determination of the composition of the nucleosynthesis products; disagreement between the amount of a substance entering the nuclear reactions, and the registered amount of a regenerated substance in the plasma bunch and in the explosion products which were deposited on accumulating screens; anomalies of the isotope composition of basic chemical elements of the material of accumulating screens; qualitative difference in the compositions of the plasma bunch and the nucleosynthesis products contained in accumulating screens; layers characterizes by the anomalous enrichment in the bulks of accumulating screens; and, finally, unidentifiable mass-peaks of atomic ions in the mass range above 220 a.m.u.

All the above-presented experimental facts were derived by different research methods and in different fields of physics, and they all contradict modern physical ideas or cannot be explained in their frameworks. However, if we assume the existence of stable superheavy chemical elements with certain properties (such as the ability to enter the nuclear reactions under the action of low-energy perturbations) which are observed, by the way, in experiments, all the contradictions disappear, and all the mentioned facts, as links, form a single chain of experimental manifestations of their existence. We are sure that the last circumstance is the most weighty argument indicating the generation of stable superheavy chemical elements in the explosive experiments.

PHYSICAL MODEL AND DISCOVERY OF SUPERHEAVY TRANSURANIUM ELEMENTS PRODUCED IN THE PROCESS OF CONTROLLED COLLAPSE

S. V. Adamenko, V. I. Vysotskii, and A. S. Adamenko

10.1. Synthesis of Superheavy Nuclei and Conditions for their Identification

During the investigations carried out from 1999 at the Electrodynamics Laboratory “Proton-21”, we have gotten a great amount of experimental data that can be interpreted as a direct or indirect confirmation of existence of the stable nuclei of superheavy elements (SHE) and superheavy nuclei (SHN).

These data can be conditionally divided into two mutually supplementing groups.

1. Data derived with the help of the “standard” spectroscopic equipment and procedures concerning the direct registration of SHN:
 - Stable isotopes with masses $270 < A < 470$ which are registered on the specimens derived in the course of experiments with various “standard” mass-spectrometers with the limiting registered values $A_{\max} = 300$ and $A_{\max} = 470$ in different laboratories, including foreign (USA, Russia) ones
 - Nonidentifiable peaks of the characteristic X-ray emission and Auger-peaks which testify to the presence of stable nuclei with $Z > 115$ to 120 and even $Z > 200$ in specimens
 - Isotopes with masses in the interval $240 < A < 4000$ which are registered by the method of Rutherford backscattering
2. Data, whose interpretation gives an indirect confirmation of the existence of superheavy nuclei:
 - Multiparticle decay of nuclear objects which is induced by a very weak perturbation and can be treated as a decay of the nuclei of SHE with $A > 1000$.

- Anomalous character of the results of the interaction of nuclear particles of the unknown nature with accumulating screens remote from the collapse point. These results can be interpreted as the interaction of superheavy nuclei with $A > 1000$, which are created in the collapse zone of a target in the process of shock compression of its central part, with the screen substance.
- Spontaneous decay of a pointlike nuclear object accompanied by the one-time emission of a very large number of fast particles. Every particle has a mass not less than the mass of an α -particle and an energy of several MeV.

Consider thoroughly some of these results which, in our opinion, convincingly testify to that the Electrodynamics Laboratory "Proton-21" has really succeeded to realize a very efficient method of creation of stable (or quasistable) superheavy nuclei and to get these nuclei in a great amount (at least 10^{13} to 10^{14} superheavy nuclei in each experiment).

First, consider briefly the main peculiarities of the performed experiments which are essential for the substantiation of the above-presented conclusion.

In the course of experiments on the realization of supercompression of the substance of a target (the anode) in a hard-current vacuum diode, we have derived the results unambiguously testifying to the formation of the collapse state in the central part of a target in each experiment. In the process of formation and confinement of this collapse, a complex of successive electron-nuclear transformations occurs, which results in the synthesis of various elements. The process of collapse of the central part of a target is accompanied by a strong heating of it. After the completion of the inertial confinement of the collapse zone, this zone is destroyed, and the plasma present in its vicinity is dispersed. This process is accompanied by the explosion-induced destruction of a part of the target. After the cessation of the collapse of the central part of a target, the synthesized elements and the plasma which are created from a target substance leave the volume, where the compression of a substance has occurred. A part of the synthesized elements is delivered to the accumulating screens surrounding the target and to the constructive elements of the internal part of the working chamber. The largest amount of synthesized elements is detected on the surface and in the volume of an accumulating screen, being at the distance of several millimeters from the collapse region.

In each experiment, we used a new (changeable) accumulating screen. But in exclusive cases, one screen was not changed during several experiments (the exclusion was made, e.g., for some screens applied later on in

the investigation of induced decays). This circumstance is related to that the partial withdrawal of a substance from the surface of an accumulating screen occurs under the action a soft X-ray emission accompanying the process of formation and destruction of the collapse zone, which is revealed in a decrease of its weight and in a change of the surface profile (the appearance of a microscopic crater of 7 to 12 mm in diameter). At the same time, we observed upon the analysis of accumulating screens that the character of the deposition of extraneous chemical elements on the accumulating screen surface testifies to that they were registered on this surface already after the change in its profile due to the withdrawal of a substance. It is obvious that the analogous withdrawal of a substance upon the repeated use of the same screens would lead to the simultaneous withdrawal of a significant part of those extraneous atoms which fell on the accumulating screen surface upon the previous shock action onto a target.

The majority of the subsequent spectrometric measurements was carried out upon the analysis of these screens. Fragments of the screens were investigated on mass-spectrometers of various types. In addition, we carried out sufficiently many studies with the remnants of exploded targets. It is worth noting that all measurements were performed in a sufficiently large time interval after the execution of specific experiments (at least in several days and more frequently in several weeks). Many measurements were carried out on the specimens which were stored after the experiment during several months. We note that no essential difference in the results of measurements was observed for different durations of the storage in all cases. This circumstance alone confirms that all registered isotopes (including superheavy nuclei) are stable or long-lived.

In the present chapter, we will carry on the analysis of conditions and results of the registration of superheavy nuclei and nuclei with anomalously large charge. The information on peculiarities and procedures of the registration of other synthesized nuclei (including the chemical elements with anomalous ratios of isotopes) and the analysis of results of this registration are given in the previous chapters of the book.

10.2. Registration of Stable Transuranium Isotopes with Standard Mass-Spectrometry Procedures

The accumulating screens and the remnants of targets taken out from the experimental setup after each shock action were studied on the equipment at the Electrodynamics Laboratory "Proton-21," at other scientific centers of Ukraine, and at several foreign laboratories. Superheavy stable nuclei were registered with the help of various mass-spectrometers.

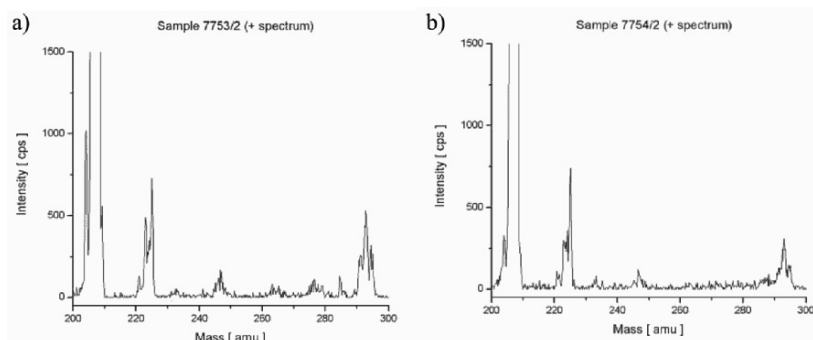


Fig. 10.1. Spectrum of isotopes on the surface of accumulating screens No. 7753 (a) and No. 7754 (b) in the region of heavy and transuranium elements $200 \leq A \leq 300$.

A part of the studies was performed at the laboratory UNITED METALS LLC (USA) with the use of an ion mass-spectrometer SIMS 030623. For this device, the range of measured masses corresponds to nuclei with mass numbers $1 \leq A \leq 300$. The primary beam was formed by Ga ions with an energy of 6 to 9 keV.

The studies were performed on specimens No. 7753 and No. 7754 which were the fragments of two different accumulating screens. The study of these specimens was realized in one month after the execution of experiments No. 4232 and No. 4233 including the shock compression of targets by a hard-current beam of electrons. The results of studies of the accumulating screens by secondary-ion mass-spectrometry at the laboratory UNITED METALS are presented in Fig. 10.1.

It is seen from the derived spectra that a great number of non-identifiable ions with mass numbers lying in the regions $A = 246$ to 248 , $A = 263$ to 265 , $A = 275$ to 278 , $A = 285$ and $A = 293$ to 297 were registered on the screens.

Analogous studies were executed in Kiev with the use of an ion mass-spectrometer SIMS of the other type, Analyzer IMS 4f (CAMECA, France). The specificity of investigations on this mass-spectrometer is related to peculiarities of the mode of scanning by the ion beam with an energy of about 10 keV on the 2D raster with the simultaneous ion etching of the analyzed region by the same beam. We have also the possibility to visualize the spatial distribution of the emission sources for specific secondary ions over this raster on a monitor display. In particular, this device admits such a mode of visualization when the raster image characterizes the spatial distribution of regions and points of the emission of any specific isotope. In addition, the spectrometer has a special operation mode “offset” which allows one to

realize the separation of monoions and molecular ions with the same summary mass number. In particular, the use of this operation mode leads to the weakening of the emission of molecular ions by at least 10 times by the practically constant emission of monoions.

These methods allow us to significantly enhance the reliability of measurements in the case where the studied isotopes are localized in the form of separate clusters, rather than are uniformly dispersed over the whole surface.

As for the heavy transuranium isotopes discovered on this mass-spectrometer, we separate the registration of ions with $A = 433$. The emission of these ions is distinguished by their unusual behavior in time. Peculiarities of their behavior were seen most clearly on the working monitor of the mass-spectrometer tuned on the registration of ions with $A = 433$. The etching of the specimen surface was carried out by Cs ions. The secondary emission of ions of the isotope with $A = 433$ occurs from a small region on the surface of one of the samples taken from the accumulating screens under study (specimen No. 7229). In the process of investigation of the specimen, it was found that the signal corresponding to the atomic ions with $A = 433$ possesses the highest intensity after the first several seconds of the etching of the studied region of the surface (this corresponds to the brightest luminescence on the device monitor). After several seconds of the etching, the luminescence brightness for this region decreases significantly [see Figs. 10.2(a), 10.3], and then the luminescence stopped, though the irradiation with Cs ions proceeded. However, after a pause in the etching of several minutes in duration, the bright luminescence restored upon the irradiation of a specimen by Cs ions [see Fig. 10.2(b), 10.3].

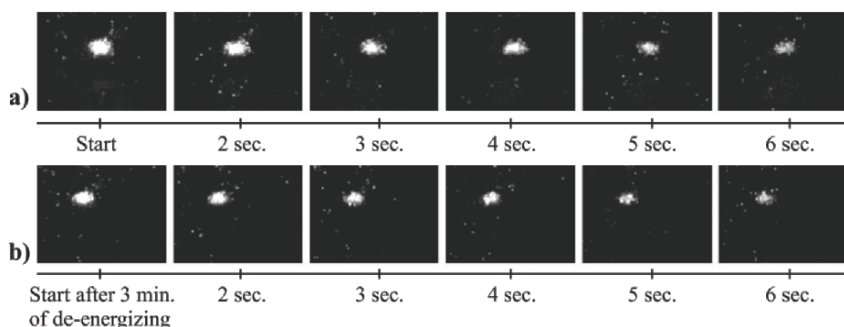


Fig. 10.2. Example of a decrease in the luminescence brightness of the spot on a mass-spectrometer SIMS display corresponding to the region of studied specimen No. 7229, on which we observed the secondary emission of ions with $A = 433$ (a), and the renewal of this brightness after a 3-min pause in the scanning of the specimen by the ion beam (b).

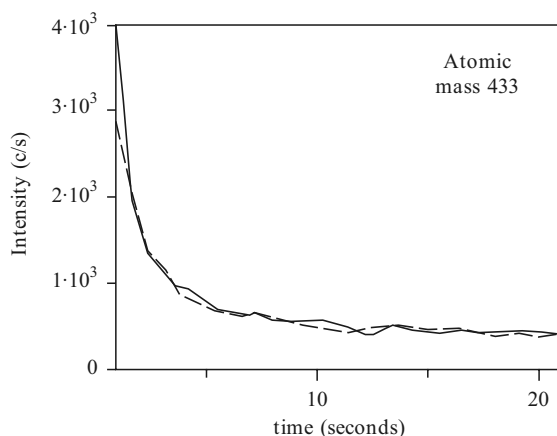


Fig. 10.3. Secondary emission intensity for ions with $A = 433$ vs the time of continuous “etching” of a specimen by the scanning beam of a mass-spectrometer SIMS.

An analogous phenomenon is surely registered also on other specimens.

For shorter intervals between series of the etching of a specimen, the analyzed region shined less intensively after the renewal of the etching. Moreover, for 3–4-sec pauses in the etching, we did not practically observe the luminescence renewal. The following dependence was clearly traced: the more the pause between the subsequent series of the etching of a specimen with Cs ions, the higher the intensity of the emission (the more the registration reliability) of the isotope with $A = 433$ after the renewal of the etching. Similar facts allow us to think about the accumulation of nuclei of the isotope with $A = 433$ on separate sections of the specimen surface. There are weighty arguments to assume that this isotope is formed due to the decay of heavier nuclei, whose masses are beyond the working range of the mass-spectrometer SIMS.

Thus, it is clear that the change in the efficiency of the emission of ions of this superheavy isotope under the action of the primary beam of Cs ions can be related to the following three mechanisms:

- the formation of nuclei of the isotope with $A = 433$ at the expense of the stimulated decay of heavier nuclei
- the decay of nuclei of the isotope with $A = 433$ at their interaction with Cs ions of the primary beam
- the fast withdrawal of ions with these nuclei at the expense of secondary-ion emission

If the rates of decay or etching are significantly more than the rate of formation of the isotope $A = 433$, then, under the action of the primary beam of Cs ions, ions of this isotope should not be registered soon after the start of the next series of etching.

The still more striking phenomena were observed on other specimen No. 7036. On one of the sections of a specimen, we found an oval region of approximately $30 \times 50 \mu\text{m}$ in size, from which the secondary emission of any ions upon the action of the primary beam of Cs ions was absent in the scope of the working range of the device, being wider than the limits of the Periodic table of elements with $1 \leq A \leq 240$. Moreover, this region of the studied specimen did not even reflect Cs ions of the primary beam, though such a reflection occurs always on any “standard” specimens! It looks a distinctive “black hole” with the absolute absorption of ions of the primary beam with an energy of about 10 keV and the absence of the emission of any secondary ions! On the outer boundary of the region, we registered a sufficiently intense emission of ions with mass numbers $A = 63$ and $A = 65$ (see Fig. 10.4). The temporal variation in the rate of emission of the indicated ions was analogous to the behavior of the isotopes with $A = 433$. During the irradiation by the primary beam, the emission of these ions stopped after some time interval. But, after the pause in the irradiation prior to the start of a next series, we registered again the intense ion current. In our opinion, these isotopes are also the products of the decay of much more heavier nuclei, whose atoms have formed this dark region. Below, we consider this mechanism in more details.

Similar dark regions with absolute absorption were also registered on some other specimens.

A great number of accumulating screens was studied at the laboratory of the scientific-industrial concern “Luch” (Russia). The studies were carried out with the use of a thermoionization mass-spectrometer TIMS “Finnigan” MAT-262. This device allows one to analyze the masses of isotopes in the mass interval $1 \leq A \leq 470$.

The data of measurements performed at this laboratory indicate the presence of superheavy atomic ions with masses in the interval $250 \leq A \leq 295$ and, in a smaller amount, nuclei in the interval $350 \leq A \leq 440$ on the studied specimens of accumulating screens. Most frequently were registered the isotopes with mass numbers: 271, 272, 277 to 280, 330, 341, 343, and 394. Some examples of the mass spectra derived in different experiments with the use of targets from different materials are given in Fig. 10.5.

As for the heaviest ions registered on this mass-spectrometer, we mention the above-discussed ions with $A = 433$ which were registered

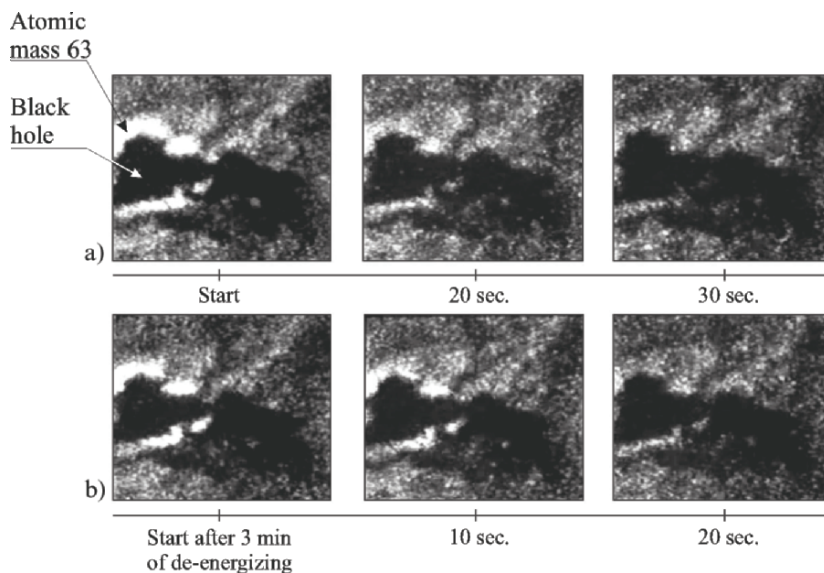


Fig. 10.4. Photos of the image of the anomalous phenomenon seen on the display of a mass-spectrometer SIMS: the practically complete absence of the secondary emission of ions in the range of masses $1 \leq A \leq 240$ from the limited area ($\sim 30 \times 50 \mu\text{m}$) of the specimen surface No. 7036. The bright luminescence on the boundary of the indicated dark zone with the emission of ions with mass numbers $A = 63$ and $A = 65$. In this case, we observed a decrease in the intensity in the course of etching of the specimen and the subsequent renewal of the emission after the 3-min pause in the scanning of the specimen by the ion beam (b).

many times with the help of the other mass-spectrometer SIMS, an analyzer IMS 4f.

The identification of isotopes was performed on the basis of the method ensuring the possibility to refer the given spectral peak to the registration of just the definite atomic ions, rather than molecular ones. To this end, we carried out the analysis of all possible combinations of atoms which would lead finally to the formation of a specific molecular ion with the prescribed mass number. In particular, the considered spectral peak did not identify with the corresponding superheavy monoatomic ion, if the spectrum of masses contained simultaneously the ions of those lighter atoms which can form the given molecule with the same total mass.

As a reference book upon the solution of the question on the identification of a spectral peak, we used the catalogs containing the maximally complete data on molecules and clusters, whose masses correspond to the

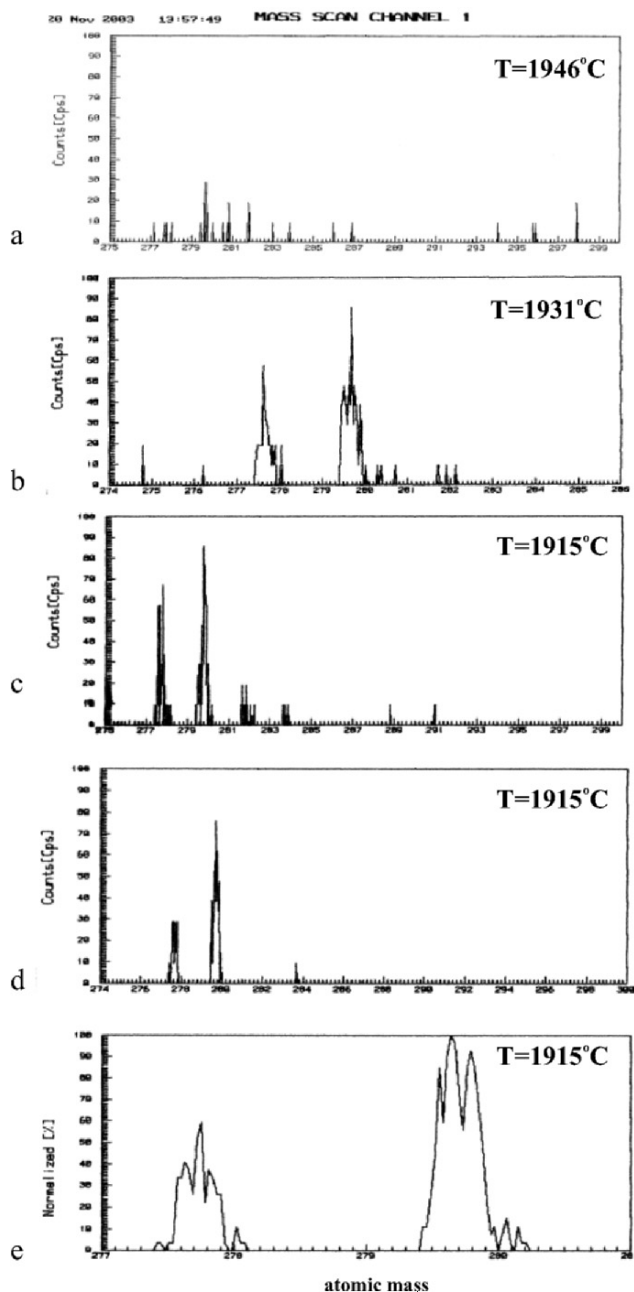


Fig. 10.5. Spectrum of isotopes in the region of great masses registered with a mass-spectrometer TIMS on the surfaces of accumulating screens corresponding to different experiments: a, b, c, d are the fragments of the panorama of spectra, e presents a more detailed view of fragment d. The temperature = 1915°C characterizes a heating element of the mass-spectrometer separated by a thin gap from a studied specimen.

working range of a device (e.g., the catalog ICPMS Interferenze Tabelle of the Finnigan MAT firm).

10.3. Identification of X-Ray and Auger Peaks of Superheavy Elements

The search for superheavy elements on the basis of X-ray spectrum microanalysis of accumulating screens was carried out on 3 devices: an X-ray microanalyzer REMMA-102 (SELMi) with an energy of accelerated electrons of 35 keV and 2 devices for X-ray fluorescence analysis. One of the latter included an X-ray tube with an accelerating voltage of 45 keV, and the other device used the γ -emission of an isotope source ^{241}Am possessing a line with an energy of 59.6 keV.

All studies were performed on specimens after a sufficiently large time interval following the execution of experiments with a shock action on a target (as a rule, not earlier than after 1–2 weeks).

The determination of the charges of nuclei was realized on the basis of the analysis of reliably registered nonidentifiable characteristic lines using the Moseley law which directly follows from the generalized Balmer formula.

Since the energy of the external action (accelerated electrons and the γ -emission of the isotope source) was significantly less than the energy of expected peaks of the characteristic emission of the K-series of superheavy elements, the identification of nuclei was carried out on the basis of the nonidentifiable characteristic emission of the L- and M-series.

The selection of a nonidentifiable peak was realized starting from the standard criterion: its amplitude should exceed the triple value of the background variance. Totally, we discovered 9 nonidentifiable peaks with energies in the range from 20 keV to 51.95 keV which can be referred to atoms with a very large charge of the nucleus.

According to the Moseley law, these peaks correspond to SHE with the following charges of the nuclei (Z): 109 to 110, 112 to 113, 115 to 117, 117 to 118, 124 to 125, 128 to 129, 141 to 143, 144 to 145, 162 to 164 if they are related to the L_α -line of the characteristic emission and 198 to 202, 204 to 208, 209 to 214, 213 to 218, 226 to 231, 234 to 239, 260 to 267, 265 to 272, 306 to 314 if they are related to the M_α -line.

The search for superheavy elements in laboratory nucleosynthesis products was realized also on an Auger microprobe JAMP-10S (JEOL, Japan). The spectra were measured at the energy of a probing electron beam equal to 10 keV.

The identification of Auger peaks was also performed on the basis of the Moseley law.

Among many nonidentifiable Auger peaks, at least three peaks can be related to nuclei with anomalously large charge. Energies of these peaks (E_e) are 172 eV, 527 eV, 1096 eV.

Depending on a specific value of the quantum number of the excited level of a bound electron in the atom of SHE, these peaks can correspond to heavy and superheavy nuclei with charge (Z): 91, 115, 140; or 113, 158, 204.

Upon a more detailed analysis of the process of measurements, we reveal anomalies in the nonstationary dynamics of the registration of non-identifiable X-ray and Auger-peaks. In particular, we found that the amplitudes of these peaks depend on the duration of the irradiation of the studied area by a probing electron beam. It was noticed that the rate of registration (the rate of collection of statistical data) of nonidentifiable peaks decreases gradually upon the long-term study of a single section of a specimen, and then the peaks cannot be registered at all. These peculiarities of the spectrometry will be considered below in the connection with the phenomenon of the induced decay of superheavy nuclei.

The procedure of these studies and the analysis of their results are described in the previous chapter in detail. There, we have also analyzed the possible sources of errors related to the process of measurements and evaluated their effect. Based on these results, we may surely consider the energies of nonidentifiable peaks and their relation to the corresponding spectral series as reliable.

Complicated is the question about the validity of using the Balmer formula and the Moseley law for the identification of superheavy nuclei by the energy of their characteristic emission and the energy of Auger-electrons.

At first glance, it seems that if the atoms of superheavy elements have the same nucleus-electron spatial structure as atoms in the scope of the known part of the Periodic table of elements, then the application of the Balmer formula and the Moseley law, which can be deduced from the Schrödinger equation, to their identification is justified.

However, it is easy to convince oneself that this question requires the additional consideration.

The matter is in that the use of the Schrödinger equation in the analysis of energy levels and in the computation of spectra of the characteristic emission and the parameters of Auger-electrons in atomic systems is grounded only in the case where the movement of atomic electrons occurs with velocities significantly less than the velocity of light. However, for the nuclei with great charge $Z \geq 100$, the mean square velocity of movement of the inner electrons $v = Z \cdot e^2/\hbar \approx Z \cdot 2.4 \cdot 10^8$ cm/s is close to the velocity of light, and we should use the Dirac equation in the computation of the

spectrum of energy levels. For the traditional representation of the energy of interaction of electrons and nuclei in the form of the Coulomb potential

$$V(r) = -Ze^2/r, \quad (10.1)$$

the Dirac equation yields that the energy spectrum of levels looks like

$$E_{nj} = m_e c^2 \left[\left(1 + \frac{\alpha^2 Z^2}{[n - (j + 1/2) + \sqrt{(j + 1/2)^2 - \alpha^2 Z^2}]^2} \right)^{-\frac{1}{2}} - 1 \right]. \quad (10.2)$$

This relation is transformed into the Balmer formula at $Z \ll 1/\alpha \approx 137$ and into the Moseley law if we take into account the electron screening of the nucleus field.

Here, $\alpha = e^2/\hbar c \approx 1/137$ is the fine structure constant, n is the principal quantum number, and $j = |l \pm 1/2|$ is the quantum number for the total momentum (the sum of the orbital and spin momenta) of an electron.

This energy spectrum has a certain peculiarity. If we take into account that the minimum value of the quantum number $j_{\min} = 1/2$ for the total momentum $\vec{J} = \vec{L} + \vec{S}$ corresponds to the zero orbital quantum number ($l = 0$) and is defined only by spin $\vec{S}$, it is clear that the formula for the energy $E_{n1/2}$ becomes complex-valued at $Z > 1/\alpha \approx 137$ and at $j_{\min}$. Such a result defines the instability threshold $Z_{cr} \approx 137$ of the Rutherford's "classical" model of atom (the system, in which one specific electron moves around the nucleus) and leads to the creation of an electron-positron pair.

If we consider the problem with a real nonpointlike (with nonzero volume) nucleus with the radius $R \approx 1.3 \cdot 10^{-13} Z^{1/3}$, this restriction is removed. But in this case, there exists the upper limit of the existence of a stable "classical" atom in the form of a planetary electron-nucleus system consisting of Z protons, $A - Z$ neutrons, and Z electrons. It corresponds to $Z_{cr} \approx 170$ (see Ref. 37). At $Z > Z_{cr}$, the one-particle Dirac equation and its solution for the electron energy levels lose sense, and it is necessary to consider a multiparticle problem with regard to the possibility of the creation of electron-positron pairs and the appearance of new energy levels of electrons which are not described by the above-presented relation.

It is necessary to note that the term "instability" is referred in this case not to the very fact of the existence of a superheavy atom, but only to the applicability of the "classical" atomic model for nuclei with very large charge. An atom proper will be stable and stationary also at $Z > Z_{cr}$. However, its energy levels will be different. The differences in the spectra will increase with the difference $Z - Z_{cr}$. The very significant variations in the characteristic spectrum are expected to be observed only in superheavy atoms with $Z - Z_{cr} > Z_{cr}$.

The analysis of the solution of the Dirac equation, E_{nj} , shows that the use of the nonrelativistic Balmer formula and the Moseley law for the determination of the charge of a superheavy nuclei with $Z < Z_{cr}$ leads to the relative error $\Delta Z/Z \sim (Z/Z_{cr})^2/n^2$.

In particular, for a pointlike nucleus, we have

$$E_{nj} \approx -\frac{Z^2 e^4 m_e}{2\hbar^2 n^2} \left[1 + \frac{\alpha^2 Z^2}{n^2} \left(\frac{n}{j + 1/2} - \frac{3}{4} \right) \right]. \quad (10.3)$$

The error will be very large upon the identification of nuclei with $Z \rightarrow Z_{cr}$ on the basis of the K_α -line (at $n = 1$), but it decreases significantly upon the use of L_α - and M_α -lines of the characteristic emission. In this case, $n = 2$ or $n = 3$ and $\Delta Z/Z \leq 10\%$ to 20% .

Starting from the above-discussed circumstances and taking into account that the above-performed identification used L_α - and M_α -lines, we should take into account that the above estimates of the charges of non-identifiable nuclei, $Z \approx 110$ to 170 , are approximate. But the difference will be small and does not exclude the fact of the registration of nuclei with very large charges. This circumstance allows us to consider that nuclei with very large Z were really registered in experiments with the use of X-ray spectrum microanalysis and electron Auger-spectroscopy, though the exact values of their charges require a correction. As for the computations giving the values of Z exceeding 170 to 200 for the identified nuclei, they should be recognized as purely evaluating. But it is obvious that such nuclei are really characterized by anomalously large charges.

It is worth noting that, for nuclei with $Z \gg Z_{cr}$, the distribution of electron levels will not correspond to the planetary electron-nucleus system even approximately. In this case, the distribution corresponds to a degenerate nonrelativistic electron gas with the approximately uniform distribution over the volume of a Wigner–Seitz cell and to a degenerate relativistic gas at $Z > \alpha^{-3/2} \approx 1700$ (see Refs. 36, 56). This question was thoroughly discussed in Sec. 11.1. For both kinds of such a gas, an analog of the characteristic emission is the transitions between the energy levels of electrons near the Fermi energy.

In particular, for a nonrelativistic gas of electrons, the minimum frequency of the transition between a pair of adjacent levels corresponds to the highest levels from the filled ones and is defined by the density of states near the Fermi energy:

$$\omega_{\min} = (1/\hbar)dE/dN = n_e^{2/3} \pi^2 \hbar / Z m_e (3\pi^2)^{1/3}. \quad (10.4)$$

For deeper levels, the frequencies of the transitions between energy levels increase gradually, and their maximum value

$$\omega_{\max} = (1/\hbar)E_F = (3\pi^2)^{2/3}n_e^{2/3}\hbar/2m_e \quad (10.5)$$

is defined by the Fermi energy.

In particular, when the density of a nonrelativistic degenerate gas is close to the limiting one, $n_e = 10^{30} \text{ cm}^{-3}$, we have

$$\hbar\omega_{\min} \approx 200/Z \text{ keV} \quad (10.6)$$

for the minimum energy of the quasicharacteristic emission corresponding to the hypothetical nucleus-electron objects with charge $Z \gg 170$. This emission corresponds to a very soft X-ray or even UV emission, whose registration is conjugated with very serious problems.

10.4. Registration of Superheavy Elements by Rutherford Backscattering

The above-presented results testify quite convincingly to the presence of stable superheavy nuclei with $A > 250$ and nuclei with very large charges $Z > 112$ in the specimens under study.

In view of the very great importance of the considered problem, we believe that the main factor defining the reliability of the results of measurements is related to the application of additional independent methods of studies and devices using basically different methods of registration. We should like to note that it is believed at present that one of the most convincing confirmations of the existence of superheavy nuclei is their direct registration by the method of Rutherford elastic backscattering. Such a method allows one to completely exclude the main factor of ambiguity of the interpretation of the mass-spectrometric experiments by separating the registrations of the individual ions of superheavy elements and the ions of molecules or molecular complexes with approximately the same total mass. Below, we consider the procedure of experiments and the results derived in more details.

10.4.1. Characteristics of Superheavy Nuclei by Rutherford Backscattering

We used the method of backward Coulomb scattering of α -particles and ions ^{14}N to study the composition of admixtures in the near-surface layers of targets (specimens). As targets, we take the accumulating screens made of chemically pure copper (99.99 %) with thickness $h = 500 \mu\text{m}$ subjected to a certain action on the above-mentioned setup (see Ref. 53).

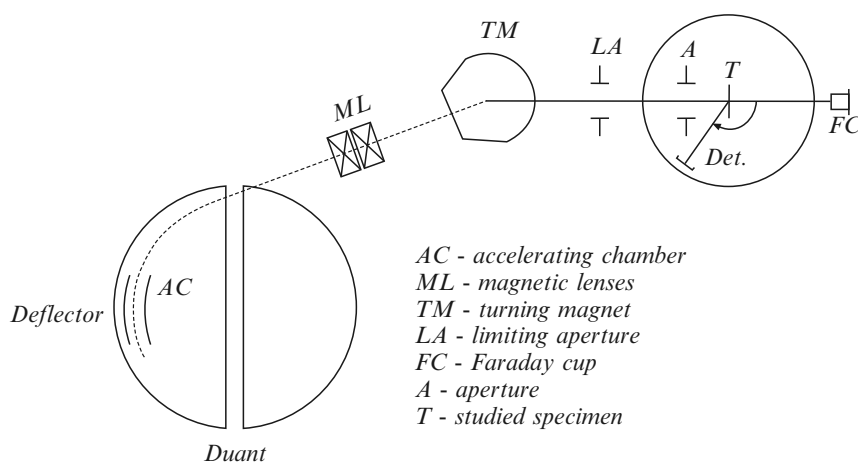


Fig. 10.6. Scheme of formation and withdrawal of the beam along the passive monochromatization channel.

In order to uniquely identify elastically scattered particles and the products of nuclear interactions, we used the $(\Delta E, E - \Delta E)$ method (see Ref. 67) for α -particles with an energy of 27.2 MeV and the time-of-flight (E, T) method for ions ^{14}N with an energy of 8.7 MeV. The energy beam dispersion was 1% of the nominal energy beam.

After the withdrawal from the accelerating chamber of a U-120 cyclotron, the beam of particles (α -particles, the current on a target was 4 nA; ions ^{14}N , the current on a target was 2 nA) was focused with magnet lenses and deflected by a turning magnet by an angle of 30° (Fig. 10.6). Then the beam was transported in a vacuumized ion-guide into the reaction room. There, after passing the limiting aperture of 3 mm in diameter, it came to the scattering chamber with the second aperture of 4 mm in diameter at its input. The studied specimen was positioned at the chamber center.

The particles scattered from a target by angles more than 90° were registered by the telescope of detectors (α -particles) or by the E detector (ions ^{14}N) (Fig. 10.7). The distance from the target to a detector was 52 cm, and the aperture on the latter was 3 mm in diameter, i.e., the solid angle $\Omega = 2.6 \times 10^{-5}$ ster. The target holder was insulated from the chamber shell. The current of particles which hit the target was supplied through a vacuum connector to a current integrator which ensured the reliable registration of the beam current in the range of 10^{-9} to 10^{-6} A.

As detectors in the $(\Delta E, E - \Delta E)$ method of identification of particles, we used surface-barrier Si detectors of 20 and 1000 μm in thickness, respectively, for the ΔE and $(E - \Delta E)$ detectors. After the preliminary and principal amplifications, signals from these detectors were supplied

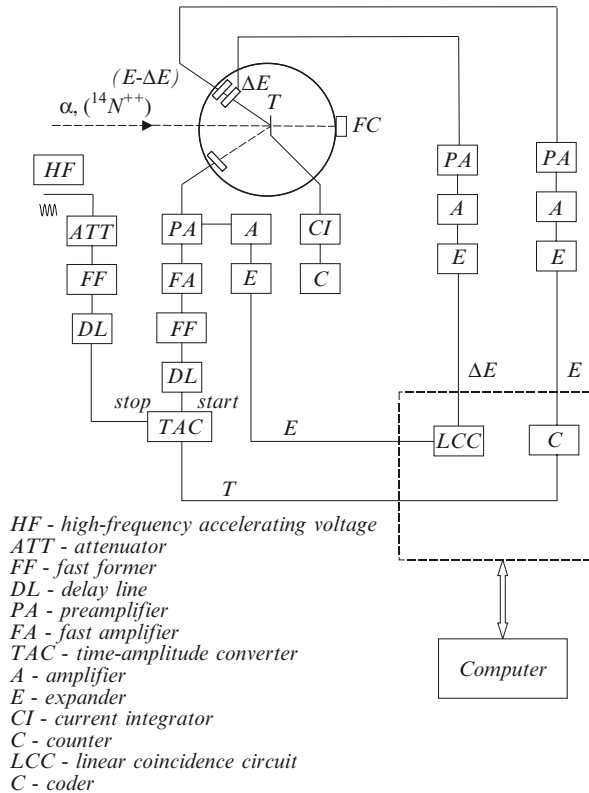


Fig. 10.7. Scheme of signal formation and registration.

through expanders to the inputs of blocks of the linear coincidence circuit and, respectively, of the coder of a 2D computer-controlled analyzer.

In the (E, T) method of identification of particles by the time of flight, we used a Si detector manufactured by a planar technology. Signals from the detector entered a preamplifier with both the fast and energy outputs. A signal from the former was supplied to the fast amplifier and, after passing a fast former, served as a “start” signal for a “time-amplitude” converter. For both the time referencing and the formation of a “stop” signal, we used high-frequency pulses of the accelerating voltage of the U-120 cyclotron. From the right duant of the accelerator, the high-frequency accelerating voltage was supplied through both a capacitance divider positioned on the shell and an attenuator to the fast former. The formed pulses came through a delay line to the “stop” input of a “time-amplitude” converter.

The “time-amplitude” converter formed analog signals with amplitude A proportional to the time of flight of a particle from the target to a detector, $T = T_{\text{start}} - (T_{\text{stop}} + T_{\text{delay}})$. In this case, the parameter T did not

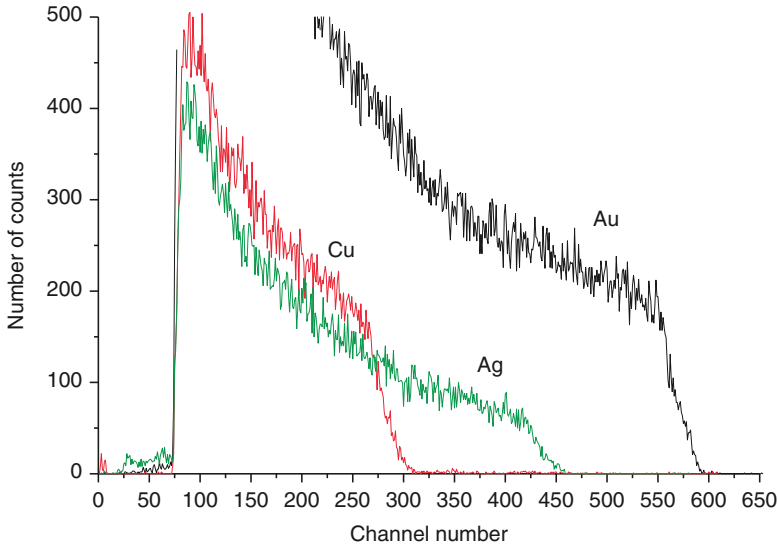


Fig. 10.8. Superposed calibrating spectra for elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 120° by specimens made of Au, Cu, and Ag.

depend on characteristics of the accelerator beam. Further, the signal T was supplied to the input of the coder of a 2D analyzer, whereas the amplified signal from the E-output of a preamplifier came to the input of the linear coincidence circuit. 2D spectra were accumulated by a MEM-3 system with the use of a spectrometric VLADO complex. We analyzed the kinematic region of elastic scattering of α -particles and ions ^{14}N exceeding the upper energy limits of elastic scattering $\text{Pb}(\alpha, \alpha) \text{Pb}_{\text{bas}}$ and $\text{Pb}(^{14}\text{N}, ^{14}\text{N}) \text{Pb}_{\text{bas}}$.

On the derived 2D spectra, we chose the loci reflecting the sort of particles incident on a specimen and constructed the energy spectra by a computer program for 2D spectra (see Ref. 68): the number of counts, from which we derived the spectra, versus the channel numbers; the number of counts versus the energy of registered particles. For the energy calibration, we took the spectra of specimens being foils of 0.5 mm in thickness made of pure (99.99 %) elements Cu, Cd, Ag, W, and Au. In Figs. 10.8 and 10.9, we give the calibrating spectra which also demonstrate the energy and mass resolutions of the method. For example, for the beams of α -particles and ions ^{14}N , the full resolutions of the setup were, respectively, 450 and 330 keV with regard to the energy dispersion ΔE of the beam in the detector volume.

A particular attention was paid to the determination of noise characteristics of the setup. To this end, we derived the spectra of the setup background (Fig. 10.10) and those of backward scattering of particles by

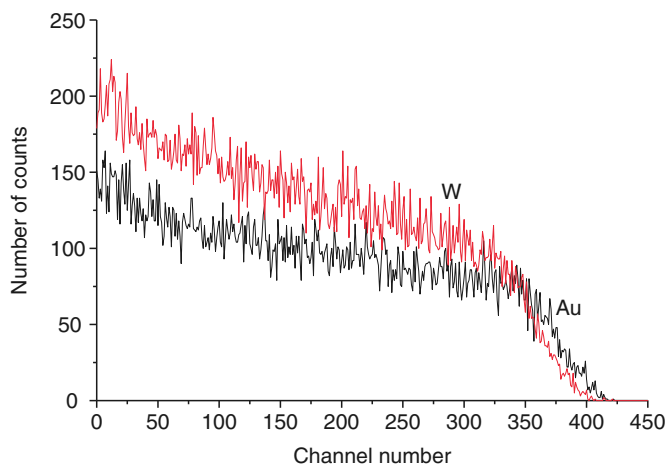


Fig. 10.9. Superposed spectra for elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 150° by specimens made of Au and W.

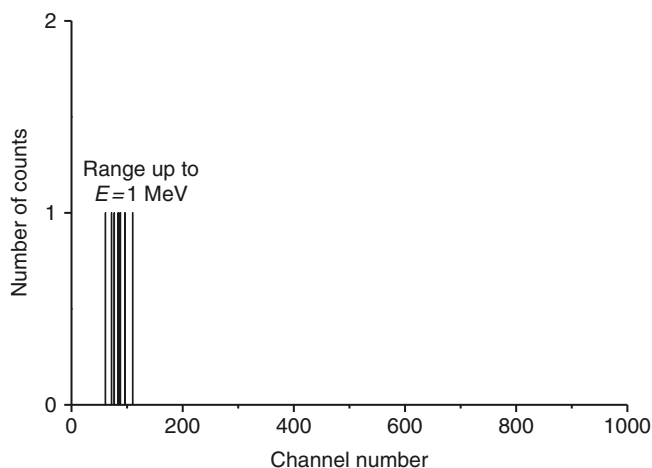


Fig. 10.10. Setup background spectrum. The exposure time equaled 6 h 10 min. The spectrometric channel was tuned for the measurement of the scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 135° . The dynamical range of an analyzer was 0 to 8.3 MeV.

specimens made of Au, Pb, and W with natural composition (Figs. 10.11–10.13) and by ^{238}Pu positioned on a dacron film with a density of about 10^{11} atom/cm ^{2} (Fig. 10.14) with the exposure time being close to the working one. As seen from the given spectra, there are no background-induced or other random events above the high-energy edge of the elastic scattering spectra from the elements composing the specimens under study for the time

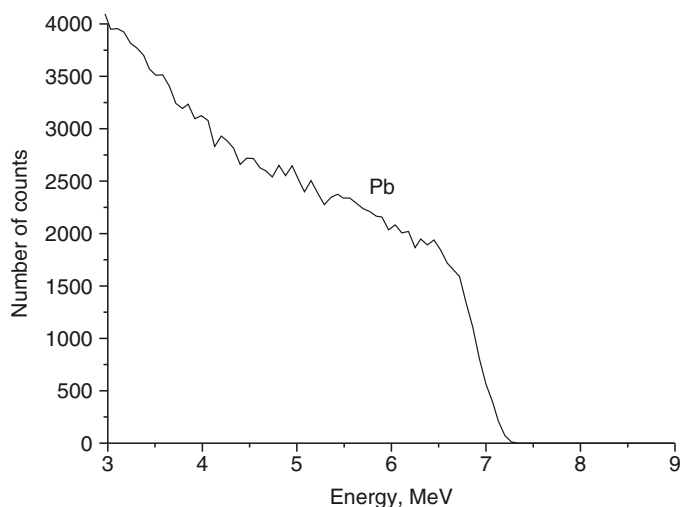


Fig. 10.11. Elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 135° by Pb. The exposure time equaled 6 h.

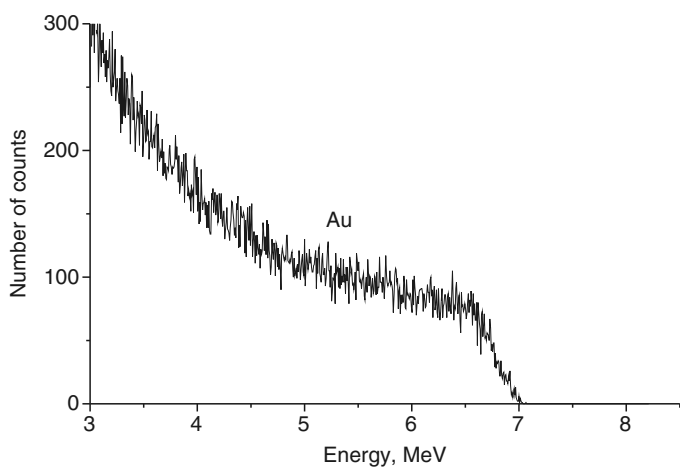


Fig. 10.12. Elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 150° by Au. The exposure time equals 4 h.

comparable with the exposure time, i.e., in the energy region where we expect to detect the events from the backward scattering of incident particles by atoms of superheavy elements.

In Figs. 10.15–10.18, we present the experimental spectra of some specimens. For specimen No.9 (the exposure time equaled 8 h 30 min),

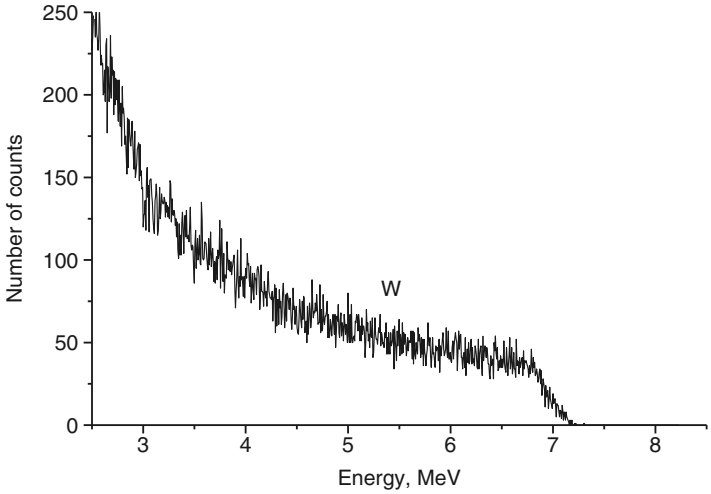


Fig. 10.13. Elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 120° by W.

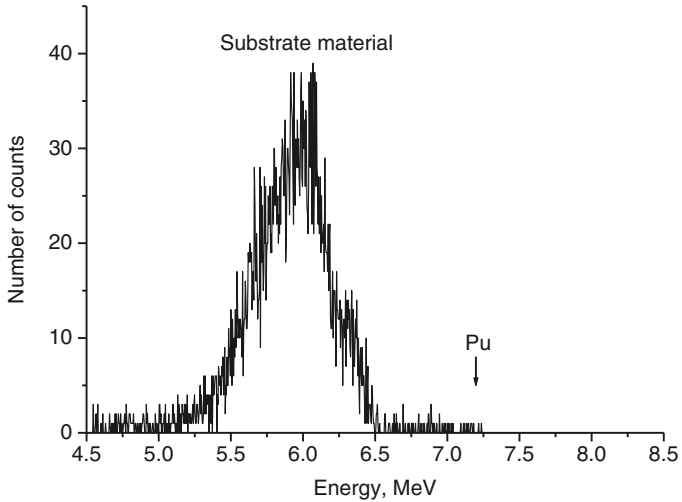


Fig. 10.14. Scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 130° . Specimen: substrate + thin film of Pu (10^{11} atom/cm ^{2}). The range of analyzed masses was 0 to 500 a.m.u.

we investigated the kinematic region of elastic scattering of α -particles by an angle of 135° , $E_{\text{kin}} = 22.0$ to 27.2 MeV. For specimens No. 8163 (the exposure time equaled 6 h 30 min) and No.37 (the exposure time equaled 10 h), the kinematic region covered 6.5 to 8.7 MeV for ions ^{14}N with the initial energy of 8.7 MeV.

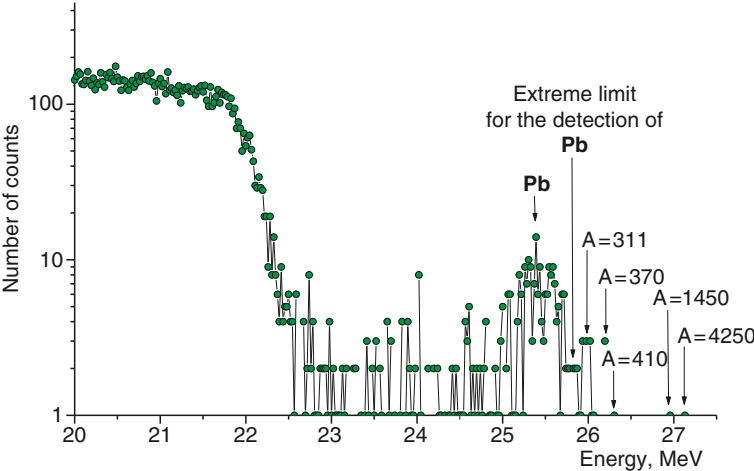


Fig. 10.15. Elastic scattering of α -particles with an energy of 27.2 MeV by an angle of 135° by specimen No. 9.

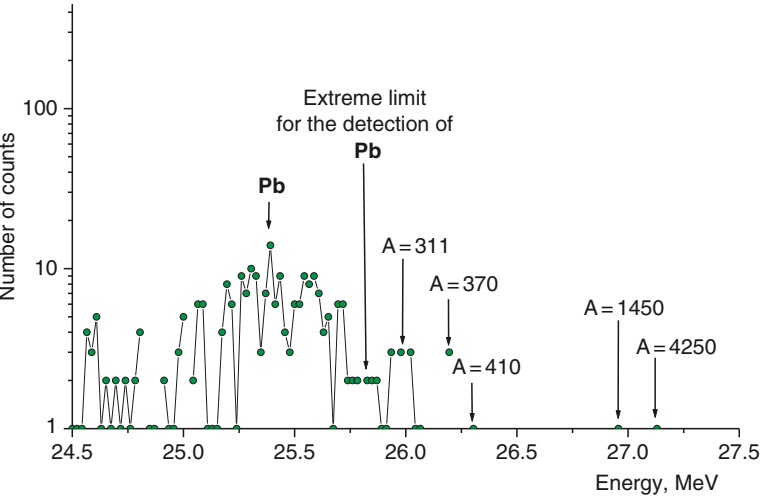


Fig. 10.16. Fragment of the spectrum of elastic scattering of α -particles with an energy of 27.2 MeV by an angle of 135° by specimen No. 9.

10.4.2. Analysis of Experimental Data

In Fig. 10.8, we give the superposed calibrating spectra of elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 120° which were derived on specimens made of Cu, Ag, and Au. While analyzing the spectra, we determined the full energy resolution of the setup which was equal to 330 keV. Figure 10.9 displays the spectra of elastic scattering of ions ^{14}N with

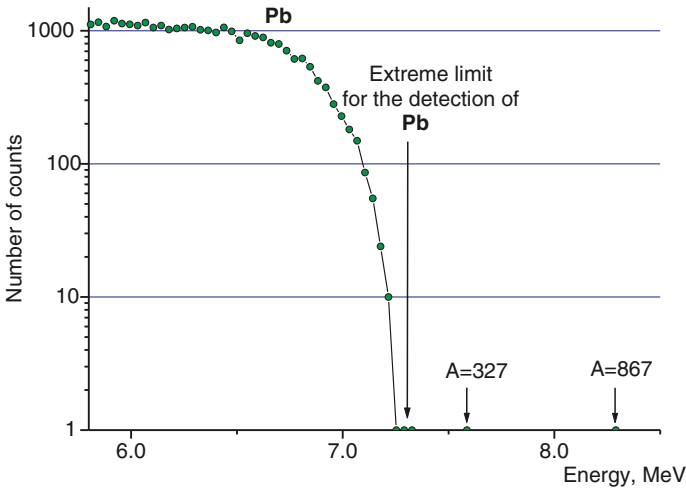


Fig. 10.17. Elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 120° by specimen No. 8163.

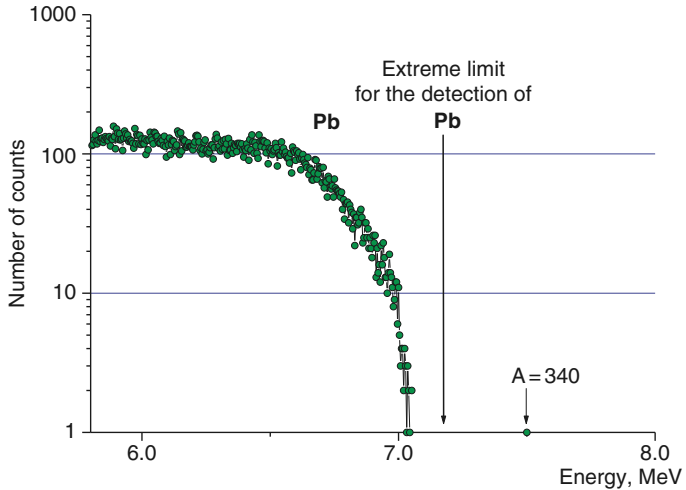


Fig. 10.18. Elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 150° by specimen No. 37.

an energy of 8.7 MeV by an angle of 150° which were derived on specimens made of Au and W. As seen, the high-energy edges of the spectra are well separated.

Figure 10.10 shows that the setup background in the range of 1 to 8.3 MeV was absent for the exposure time of 6 h 10 min, which corresponds to the dynamical range of detected masses of 25 to 900 a.m.u. We also did not observe any setup background for the exposure time of 6 h

under elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 135° by the Pb specimen (the dynamical range was up to 900 a.m.u.) (see Fig. 10.11) and by an angle of 150° for 4 h by Au in the dynamical range up to 700 a.m.u. (Fig. 10.12). The difference in the dynamical ranges is related to the tuning parameters of electronic blocks of the spectrometric channel. In Fig. 10.14, we give the spectrum of elastic scattering of ions ^{14}N with an energy of 8.7 MeV by an angle of 130° which was derived on the specimen being a dacron film with introduced ^{238}Pu , whose amount was 10^{11} atom/cm ^{2} by the estimates based on the specimen radioactivity. The exposure time equaled 7 h. This figure along with Figs. 10.10–10.12 shows the possibility to register heavy elements in a light matrix at the level of 10^{11} atom/cm ^{2} .

It is seen from the spectra presented in Figs. 10.15–10.18 that the dynamical energy range corresponding to the elastic scattering of incident particles by the nuclei of atoms with masses of above 260 a.m.u. includes the events which can be referred to the processes of elastic scattering of α -particles and ions ^{14}N by the nuclei of superheavy elements with $A = 310$ to 4500.

Table 10.1 contains the resultant data on the elastic scattering of α -particles and ions $^{14}\text{N}^{++}$ by the accumulating screens in the energy ranges $E_\alpha = 26.00 \pm 0.45$ to 27.20 ± 0.45 MeV and $E_N = 6.5 \pm 0.3$ to 8.7 ± 0.3 MeV,

Table 10.1. Parameters of the registration of elastically scattered α -particles and ions $^{14}\text{N}^{++}$ by specimens.

Specimen	Type of particles, scat. angle	Energy of particles (ions), MeV		a.m.u.	Charge of a nucleus
		Accelerated	Scattered		
No. 9 Cu	α -particles 135°	27.2	26.0	311_{-86}^{+190}	114
			26.2	370_{-120}^{+280}	134
			26.3	410_{-140}^{+420}	146
			26.9	$1450_{-930}^{+\infty}$	372
			27.1	$4250_{-3570}^{+\infty}$	712
No. 8163 Cu	$^{14}\text{N}^{++}$ 120°	8.7	7.6	327_{-80}^{+145}	122
			8.3	$865_{-395}^{+\infty}$	259
No. 37 Cu	$^{14}\text{N}^{++}$ 150°	8.7	7.5	340_{-75}^{+130}	126

respectively. Errors in the determination of masses were calculated from the total dispersion of the beam energy, and the nucleus charge was evaluated in the frame of the liquid-drop model by the Weizsäcker formula for stable nuclei:

$$Z = A / (1.98 + 0.015A^{2/3}). \quad (10.7)$$

Table 10.2 contains the statistical information about the events corresponding to the elastic scattering of α -particles with the energy $E_\alpha = 27.2$ MeV by the nuclei of SHE with masses $296 \leq A \leq 310$ which were located on the accumulating screen surface. On the whole, during the execution of 8 experiments, we got 72 events which can be interpreted as the registration of superheavy nuclei with masses in the interval $296 \leq A \leq 310$.

Summarizing, we note that the method of Rutherford backscattering has confirmed the existence of stable superheavy nuclei and allowed us to determine their masses. During the experiments, we registered a great number of the events of backscattering related to the nuclei with masses in the interval $A = 296$ to 310 and also to the nuclei with $A \approx 311, 327, 340, 370, 410, 865, 1250, 4250$.

Table 10.2. Characteristics of the registration of superheavy nuclei with masses $296 \leq A \leq 310$ in the series of experiments on the elastic scattering of α -particles with the energy $E_\alpha = 27.2$ MeV by the accumulating screen surface.

Series	A	E , MeV	σ^* , nb	$\sigma_{\text{background}}$, nb	$P_1^{**} \times 10^{13}$	$N^\dagger$
1	296 to 310	25.5 to 26.5	8 to 10	0.1	5.4	9
2	296 to 310	25.5 to 26.5	14	0.1	9.6	16
3	296 to 310	25.5 to 26.5	12	0.1	6.6	11
4	296 to 310	25.5 to 26.5	4	0.1	3.0	5
5	296 to 310	25.5 to 26.5	13	0.1	8.4	14
6	296 to 310	25.5 to 26.5	12	0.1	6.6	11
7	296 to 310	25.5 to 26.5	2	0.1	0.6	1
8	296 to 310	25.5 to 26.5	4	0.1	3.0	5

* total cross section of backscattering

** probability of the appearance of one event in a physical region above the kinematic limit

† number of events corresponding to the registration of SHN

A part of the derived values of A corresponds to (with regard to errors of the method of backscattering) those mass numbers which were derived by “standard” mass-spectrometry with the use of TIMS and SIMS spectrometers and were considered above.

Moreover, in view of the statistics, the duration of measurements, the backscattering angle, and the solid angle of a registration, we can assert that the surface density of each discovered type of superheavy nuclei located in the thin surface layer of the studied accumulating screens of about $100\rho A$ in thickness is $\sigma \approx 10^{11}$ to 10^{12} cm^{-2} .

10.5. Induced Decay of Superheavy Nuclei with the Help of a Beam of Oxygen Ions and Upon the Action of Laser Emission

Upon the investigation of accumulating screens by nondestructive analytic methods, in particular by X-ray spectrum microanalysis and Auger-spectrometry, nonidentifiable peaks with X-ray emission were reliably registered in spectra. On the basis of the general systematization of the energy levels of electrons in atoms, we concluded that these peaks can be associated with the characteristic emission of the electron system of SHE. This question was thoroughly considered in the previous chapter.

Upon the more detailed study of these peaks, we discovered some anomalies. In particular, we observed that their amplitude decreases with increase in the duration of the irradiation of a studied section by a probing electron beam. It was also noticed upon the study of specimens by X-ray spectrum microanalysis that the decrease in the intensity of peaks is greater and occur more frequently, than that upon the study by Auger-spectrometry.

It is known that both analytic methods are nondestructive, and their application does not lead to structural changes in the target and to the emission of secondary ions. The discovery of the influence of the irradiation duration and the electron beam energy on the intensity of spectra allows us to assume that the probing irradiation induces a very essential transformation of the internal structure and charge state of the nuclei of those elements, whose spectrum contains nonidentifiable peaks. In particular, we put forward a hypothesis that, in this case, we observed the induced decay of the nuclei of SHE under the weak (on the scale of nuclear processes) external action of the electron beam.

Such an explanation of the registered effect agrees well with peculiarities of two considered types of measurements. Since the energy of electrons of the primary electron beam upon X-ray spectrum microanalysis is 3–4 times more than the energy of electrons upon Auger-spectral analysis, it is logical that the decay of nuclei can be easier induced in the first case.

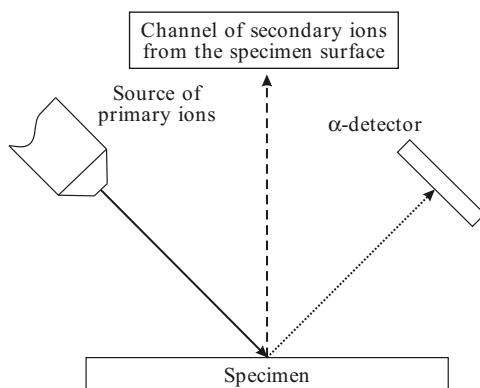


Fig. 10.19. Scheme of the position of a Si detector of α -particles in the vacuum chamber of a secondary-ion mass-spectrometer SIMS in the experiments on the induced decay of superheavy nuclei.

In order to check this assumption and to register products of the induced decay of superheavy nuclei by using a secondary-ion mass-spectrometer Cameca IMS 3F, we carried out the experiment with the use of an ion microprobe. We intended to provoke the fission of the nuclei of SHE by the irradiation of the surfaces of accumulating screens by accelerated ions of oxygen.

The experiment scheme is shown in Fig. 10.19. The accumulating screens, on which SHE were probably located, were mounted on the working plate of the device. An Si detector of α -particles was installed directly in the vacuum chamber of a mass-spectrometer. The active part of the detector was near the optical axis of the secondary ion channel. The distance between a specimen and the detector was approximately 3 cm. Specimens were irradiated by the beam of ions O_2^+ with an energy ~ 12.5 keV at a pressure of residual gases in the working chamber $\sim 5 \cdot 10^{-9}$ Torr. The current of primary ions was in the range 1 to $5 \mu\text{A}$. For measurements, we used a specially designed system including the detector of α -particles, preamplifier, final amplifier, counter of pulses, and digital oscillograph for the visual estimation of the pulse amplitudes.

The workability of the measuring scheme was verified with the use of a source of α -particles. We also established experimentally by using the testing β - and γ -active sources that the measuring scheme of the experimental setup does not register β - and γ -emission (at least in the range up to 3 MeV), which allowed us to neglect such a source of errors in what follows. We carried out also the registration of the background α -emission upon the switched-off primary beam and without a specimen and the registration of events

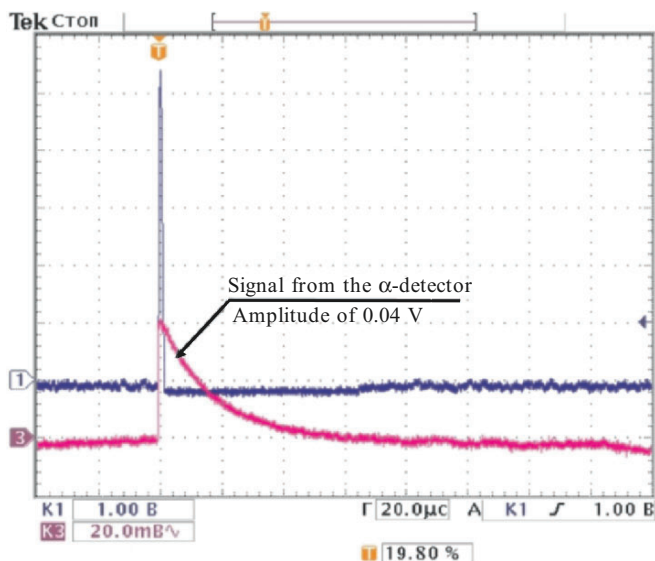


Fig. 10.20. Pulse from a background α -particle registered in the study of a specimen not participating in the experiment.

upon the etching of the control specimen made of Cu (not participating in the experiment) by the primary ion beam. In this case, we registered one event for 2 h on the average. The amplitudes of pulses from the background α -emission, which were registered by the detector, varied in the range of 40 to 70 mV. The example of a pulse from a background α -particle is shown in Fig. 10.20.

The α -detector signal amplitude registered by an oscillograph is proportional to the energy and charge of a registered particle. Since the background α -particles emitted upon the decay of the U-Ra- and Th-series have energies from 5.5 to 8.8 MeV, the change of the energy of a registered α -particle by 1 MeV corresponds to a change of the pulse amplitude by 7 to 8 mV. Upon the calibration of the measuring scheme with the help of a Np source emitting α -particles with an energy of 4.8 MeV, the signal amplitudes did not exceed 35 mV.

The testing of the method of studies began on specimen No. 53. This specimen was used as an accumulating screen in 6 subsequent experiments upon the shock action on lead targets. First, upon the etching of a specimen by the beam of ions, we used a raster of $250 \times 250 \mu\text{m}$ in size. In this case, we registered no anomalous signals. Then, with the purpose to increase the probability of the interaction of a primary beam with the nuclei of SHE, the studies were carried on by an ion probe. In this case, the etching region diameter was approximately $100 \mu\text{m}$. Such a decrease of the etching region

would lead to an increase in the probability of the interaction of ions of the primary beam with the nuclei of SHE and simultaneously to a decrease in the probability for the ion probe current to fall on sections of the surface where the nuclei of SHE can be localized.

Each section of the accumulating screen was etched with a probe for about 1 h. If no signals were registered for this time, excluding for the rare signals of background α -particles, the probe was moved to an adjacent point of the specimen.

On one of the sections of this specimen in 25 min of the etching, we began to register the pulses, whose amplitude exceeded the amplitude of background events by a factor of several tens. During the first and second hours, we registered, respectively, 7 and 4 pulses. Then the registration of pulses stopped. On two other sections, we registered, respectively, 3 and 4 pulses during the first hour of the etching. The registered pulses were similar to the pulses from background α -particles by the form and considerably exceeded them by the amplitude. The amplitude of the former varied from 0.1 to 4 V. Under the condition of the isotropic emission of particles with regard to the solid angle, in which the detector is placed relative to the specimen, the total number of particles emitted from the accumulating screen was about 500. In Fig. 10.21, we show the oscillograms of two pulses registered on specimen No. 53. The amplitudes of these pulses testify to the registration of α -particles with energies of hundreds of MeV or to the registration of heavier products of the decay of nuclei. On the rest sections of specimen No. 53, we found no pulses for the 3 h etching.

A somewhat other pattern was derived upon the study of specimen No. 69. This specimen was also used as an accumulating screen in 6 subsequent experiments upon the shock action on lead targets. On one of the sections of the specimen in 5 min of the etching, we registered 2 series of events (110 and 30 pulses, respectively). The interval between series was

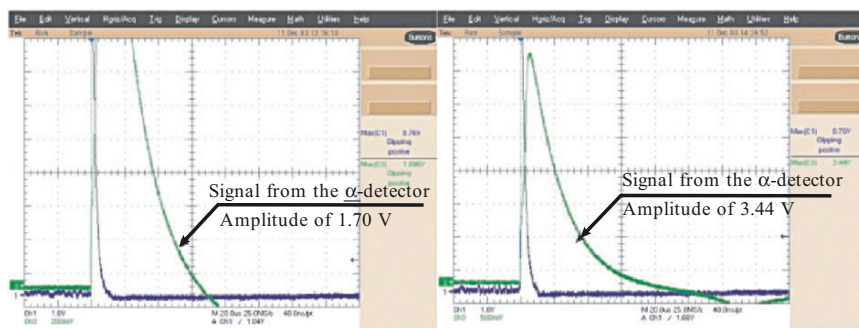


Fig. 10.21. Pulses registered upon the study of specimen No. 53.

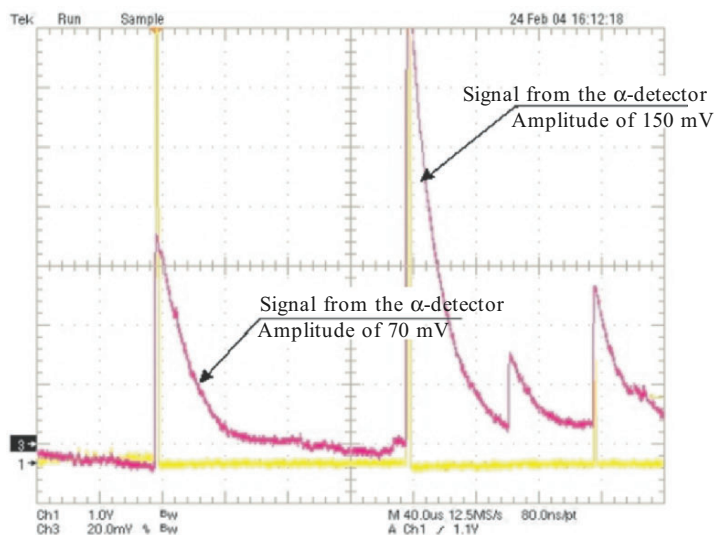


Fig. 10.22. Fragment of the series of events registered upon the study of specimen No. 73.

approximately 1 min. The duration of each series was several seconds. But we failed to register the amplitudes of these pulses due to the absence of the oscillograph during the execution of the experiment which was due to some technical reasons.

However, a similar phenomenon was registered also upon the study of specimen No. 73 (the specimen was used as an accumulating screen in 3 experiments on lead targets). In 2 to 3 min after the start of the etching of one of the sections of the specimen under the current of primary ions of $5\mu\text{A}$, we registered 257 pulses during approximately 1 s. In Fig. 10.22, we demonstrate the oscillogram of a fragment of the event. This event was also registered by the other α -detector, whose calibration indicated that the change in the energy of a registered α -particle by 1 MeV corresponds to the change in the pulse amplitude by 5 mV. Under the condition of the isotropic emission of particles with regard to a solid angle, in which the detector was relative to the specimen, the rated number of particles emitted from the accumulating screen exceeded 1500. The horizontal scanning of the oscillograph can contain only 4 pulses. The interval between these pulses was from 60 to $100\mu\text{s}$. The amplitudes of two first ones (from 4 registered pulses) were, respectively, 70 and 150 mV, which exceeds considerably the amplitudes of background signals and corresponds to energies of α -particles of 14 and 30 MeV.

Besides the above-described experimental results, upon the action of a primary beam on some sections of the specimens, we registered some

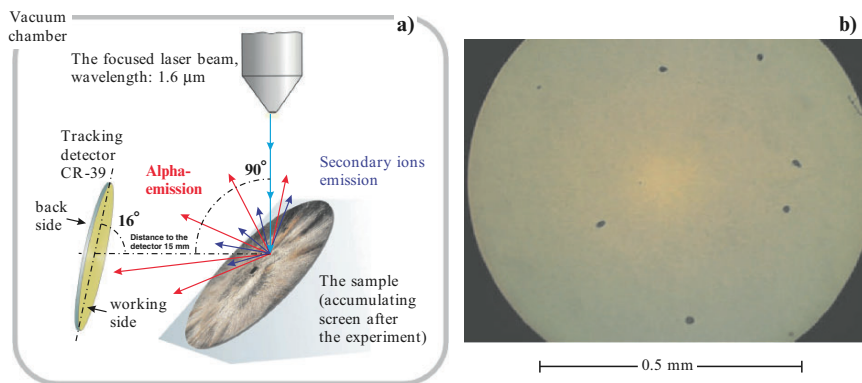


Fig. 10.23. Scheme of the experiment on the study of the induced decay of SHE with the help of laser emission (a); a fragment of the surface of a track detector after the series of laser pulses and the subsequent etching of the detector surface (b).

excess of the α -background (about 2 particles for 1 h at a background level of 0.5 particle/hour) and small series of events from 4 to 25 pulses for several seconds many times. In this case, the amplitudes of some signals exceeded those from background α -particles by a factor of 1.5 to 2.

We carried out also the experiments on inducing the decay of SHE with the use of a laser emission (Fig. 10.23).

Fragment of the working surface of the detector (after the etching) with tracks of the 5 to 8-MeV α -particles. The total number of tracks on the working surface of the detector is 149. The number of tracks on the back side (background events) is 18. The etching duration is 6 h in 25 % NaOH at 70°C.

The experiment scheme was analogous to the scheme of the etching of a specimen by ions of oxygen. Instead of a source of ions, we used a focused laser beam. The registration was realized with help of a track detector CR-39. The experiment was carried out in vacuum at a residual pressure of $1 \cdot 10^{-4}$ Torr. The laser emission energy per pulse was about 50 mJ, pulse duration was 14 ns, emission power density on a target was $2 \cdot 10^9$ W/cm², and wavelength was 1.06 μ m. On specimen No. 10344, 50 points were subjected to the action of a laser beam. At each point, we directed 10 pulses. All this time, the track detector was in the vacuum chamber at a distance of about 15 mm from the specimen. The detector was fixed so that its back side was accessible to the action of the residual atmosphere of the vacuum chamber, but α -particles emitted by the surface of the studied accumulating screen cannot come to it. Hence, the back side of the detector can be considered a reference detector registering the background of α -emission into the vacuum chamber

during the execution of the experiment. After the etching, we discovered about 150 tracks of α -particles with energies up to 8 to 10 MeV on the face (working) surface of a track detector. Simultaneously, we registered only 18 tracks of α -particles with energies up to 8 to 10 MeV on the back side of detectors.

Such great difference of the total numbers of signalling and background tracks, in our opinion, testifies unambiguously to the induced character of the decay of nuclei under the action of a laser irradiation.

Thus, on the basis of the derived experimental data, we may draw the following conclusions:

- The action of the beam of oxygen ions on inhomogeneous microporations of a substance, which contains the explosion products of targets and is located in the surface layer of an accumulating screen, initiates the processes that can be related, with a high probability, with the decay or fission of the nuclei of SHE present in the explosion products of a solid target.
- We observed two different scenarios of the decay: fast and slow ones; in the first case, several hundreds of particles, for which the product of energy by charge is several times more than that of background α -particles are emitted for several seconds; in the slow decay, the process is running for 1 to 2 h and is accompanied by the emission of several tens of particles, for which the product of energy by charge is several tens of times more than that of background α -particles.
- The action of pulses of the optical laser emission with low energy on the surface of an accumulating screen stimulates the emission of α -particles with energies up to 8 to 10 MeV; the mean efficiency of such stimulating action upon a “blind” (i.e., random and unsighting) action on a target is small and is equal approximately to two α -particles per laser pulse.

10.6. Induced Decay of Superheavy Nuclei with the Help of a Beam of Cu Ions

The independent series of studies of the process of induced decay of non-identifiable superheavy nuclei was carried out on the specially designed setup with the use of accelerated Cu^{++} ions.

The fission of SHN can be the most reliable proof of their existence and their unusual shape. This phenomenon possesses a number of kinematic peculiarities which allow one to unambiguously plan the experiment and admit no ambiguity in the interpretation of its results. Therefore, by assuming a low binding energy of SHN, we carried out a series of experiments on

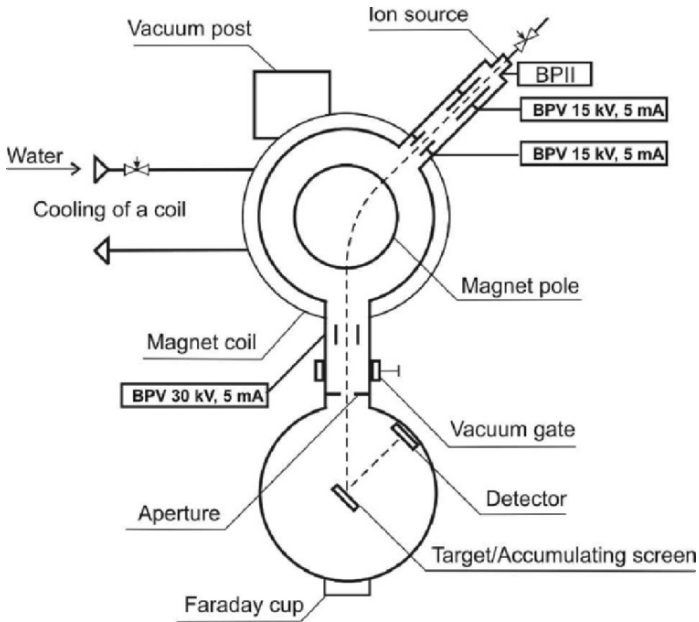


Fig. 10.24. Block diagram of the experimental setup.

the excitation of decay modes of the metastable states of such nuclei. For the excitation, we used beams of $^{65}\text{Cu}^{++}$ with $E = 20 \text{ keV}$. Such beams were formed with the source of heavy ions by the scheme given in Fig. 10.24.

Main units of the setup:

- source of ions $^{63,65}\text{Cu}^{++}$; $E = 10 \text{ to } 20 \text{ keV}$; $I = 5 \mu\text{A}$
- channel of formation of the beam of $^{63,65}\text{Cu}^{++}$
- channel of separation of ions $^{63}\text{Cu}^{++}$ and $^{65}\text{Cu}^{++}$
- route of formation of the beam of $^{65}\text{Cu}^{++}$
- target in the reaction chamber, where the beam of $^{65}\text{Cu}^{++}$ with $E = 10 \text{ to } 20 \text{ keV}$ interacts with the surface of an accumulating screen
- system of accumulation and processing of the data derived during the registration of products of the interaction of doubly charged ions with the surface of accumulating screens

On the interaction of a doubly charged ion ^{65}Cu with $E = 10 \text{ to } 20 \text{ keV}$ with atoms of the material of accumulating screens, a number of elastic and inelastic processes occurs.

In the process of deceleration related to elastic collisions, $^{65}\text{Cu}^{++} + ^{63,65}\text{Cu}^{++}$, $^{65}\text{Cu}^{++} + \alpha$, the energy of an incident ion is transferred to atoms of the material of the accumulating screen without change in their energy and charge states. These processes have discrete character. Such a nuclear scattering is realized into vacuum or into depth of the target.

Upon the bombardment of the surface of metals, the sort of an incident particle (charged or neutral) is of no importance since the incident ion neutralizes its charge on the way of several nanometers for the time close to that of orbital motion of an electron (about 10^{-16} s).

In the description of the interaction of an incident ion with atoms of a target, we neglect the influence of adjacent atoms.

The process of elastic collisions of ions at $E = 10$ to 20 keV includes:

- reflection of ions into vacuum
- dispersion of atoms of the material
- heating of a target
- formation of defects
- implantation of ions ^{65}Cu into the material of a target-accumulating screen
- channeling

The processes of inelastic scattering of ions $^{65}\text{Cu}^{++}$ with $E = 10$ to 20 keV induce the excitation of the electron subsystem of atoms of the accumulating screen.

Relaxation channels:

1. ion-photon emission
2. ion-electron emission
 - (a) ion-stimulated desorption
 - (b) change in charge states of atoms of the accumulating screen material
 - (c) formation of radiative effects

We assume that the emitted ions move from the surface in depth of the accumulating screen, being neutral or quasineutral with holes in deep electron shells which are screened by outer electrons.

On the impact of an ion $^{65}\text{Cu}^{++}$ with $E = 10$ to 20 keV, significant disturbances occur in the accumulating screen.

The energy brought by an ion $^{65}\text{Cu}^{++}$ influences a very small volume. At $E = 10$ to 20 keV, a movement of the ion in the Cu accumulating screen along the distance of 8×10^{-9} m is associated with a nonequilibrium process accompanied by high-velocity heating of the decelerating layer (about 10^{14} K/s).

This leads to the ionic etching of the surface of the accumulating screen and to the knocking out of 7 to 8 atoms of the accumulating screen material with $A = 40$ to 70 a.m.u. by every ion ^{65}Cu . The accumulating screen surface changes its structure and element composition.

In the near-surface layers of targets, the electron-involved processes between $^{65}\text{Cu}^{++}$ and atoms on the accumulating screen surface with $A > 200$ ($Z > 80$) run. The following processes are possible:

1. resonance charge transfer in the processes of adiabatic tunnelling
2. charge transfer in the Auger-processes
3. nonadiabatic tunnelling

In experiments, we used a low-noise α -spectrometer. A semiconductor Si-detector registered charged particles, products of the induced fission, in inverse geometry. For the beam current of $5\ \mu\text{A}$ flowing through the accumulating screen-target containing SHN, the loads were at most 5×10^2 particle/s. For such loads, the spectrometer of charged particles worked in the mode of self-coincidences and the time rejection of superpositions and allowed us to reach the signal-to-noise ratio to be about 10^{15} . The gaging spectra with sources ^{226}Ra and $^{238,239,240,242}\text{Pu}$ and with different time exposures (up to 100 h of the continuous accumulation of events) are given in Figs. 10.25–10.27.

It is necessary to briefly dwell on the procedure of measurement of α -spectra and the gaging of the α -spectrometer. Si-detectors which were used on the registration of the possible decay products of SHN have a linear response function for elastically scattered ions $^{65}\text{Cu}^{++}$ with

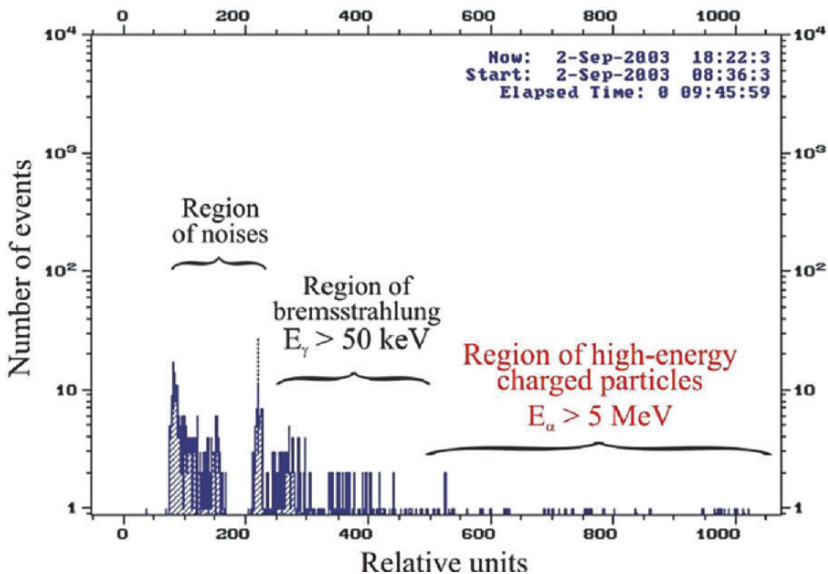


Fig. 10.25. Bremsstrahlung spectrum showing the distribution over channels (discrete decay states) of γ -quanta and α -particles.

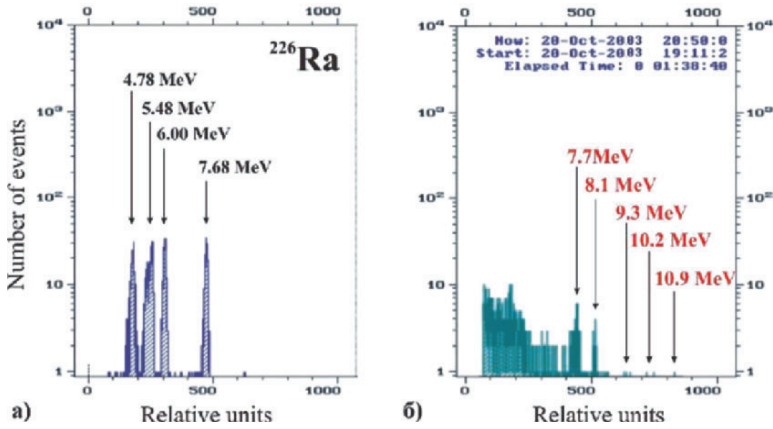


Fig. 10.26. Energy spectrum of α -particles. (a) a gaging source with ^{226}Ra , (b) α -particles from specimen No. 37.

$E_{\text{Cu}} = 10$ to 20 keV, γ -quanta of various nature with $E_{\gamma} = 2$ to 70 keV, and α -particles with $E_{\alpha} = 0.1$ to 30 MeV. It is natural that the response functions for such distinct products of the interaction of low-energy heavy ions with the accumulating screen surface lie in different dynamic ranges of the α -spectrometer.

Even in the mode of 1D analyzer, this allows us to reliably separate the processes $A(^{65}\text{Cu}, ^{65}\text{Cu})$;

$$A + ^{65}\text{Cu} \rightarrow \gamma, \text{ and } A + ^{65}\text{Cu} \rightarrow B^* \left\{ \begin{array}{l} \rightarrow C \\ \rightarrow \alpha \\ \rightarrow \gamma \end{array} \right.$$

Here, A is nucleus-target, ^{65}Cu is incident Cu ion, B^* is formed compound system, C is nucleus-residue derived after the decay of the compound system.

These dynamic ranges of the α -spectrometer are well seen on Fig. 10.25 derived in the on-line mode.

The spectra of induced fission of SHN derived in the on-line mode with the $^{65}\text{Cu}^{++}$ beam at $E = 20$ keV show 3 regions of the dynamic range of registered events (Fig. 10.26):

1. the region of registration of X-ray and γ -emission in the range 2 to 60 keV (γ -quanta corresponding by energy to a Si-detector with $h = 300 \mu\text{m}$)
2. the region of registration of α -particles with $E_{\alpha} = 3$ to 7.6 MeV
3. the region of registration of α -particles with $E_{\alpha} = 8$ to 15 MeV

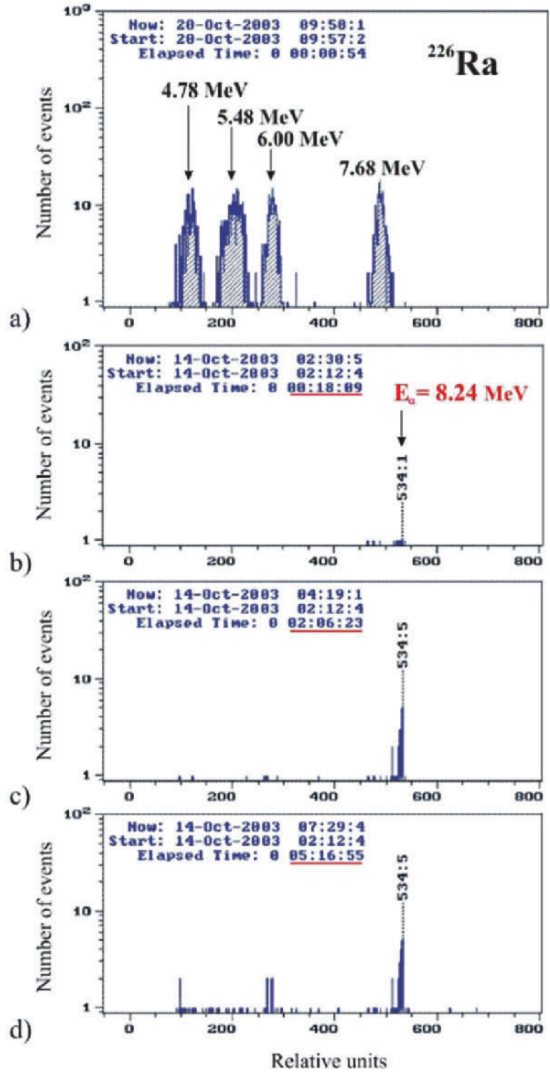


Fig. 10.27. Energy spectrum of α -particles (a) a gaging source with ^{226}Ra . The total dynamics of accumulation of discrete α -lines from specimen No. 37, the accumulation time: (b) 18 min, (c) 2 h 6 min, (d) 5 h 16 min.

We can say about a significant statistics only for the first two regions.

In the region $E_\alpha = 8$ to 15 MeV , the used system of identification of the energy and types of particles does not give us sufficient grounds to determine something in a single-valued manner due to both the small statistics and the spectral range of unit events almost continuous in energy. We can only assert that the energy range of registered events reaches $E_\alpha = 25 \text{ MeV}$.

The second series of experiments on the induced decay of SHN was performed by the off-line scheme. The accumulating screens were irradiated by low-energy beams of $^{65}\text{Cu}^{++}$ with $E = 20\text{ keV}$ for 24 h. After the irradiation, the accumulating screens were positioned in a low-noise α -spectrometer, and then we accumulated events forming the energy spectrum of α -particles. In Fig. 10.27, we show the total dynamics of accumulation of the discrete α -lines corresponding to α -decays of parent SHN.

In Figs. 10.26–10.31, we present the energy spectra of α -particles from specimens Nos. 37, 8384, 8381, 50, 8383 after the action of the beam of $^{65}\text{Cu}^{++}$ with $E = 20\text{ keV}$ (see Table 10.3).

Specimens in the amount of 15 units were supplied for experiments by the Electrodynamics Laboratory “Proton-21” specially for the search for SHN in these specimens, because the presence of SHN in them was indicated by the previous mass-spectrometric studies.

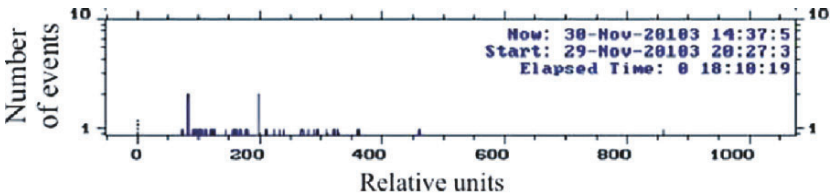


Fig. 10.28. Energy spectrum of α -particles from specimen No. 8384 after the Cu^{++} action.

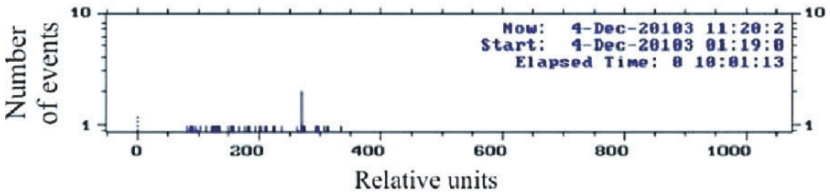


Fig. 10.29. Energy spectrum of α -particles from specimen No. 8381 after the action of $^{65}\text{Cu}^{++}$ with $E = 20\text{ keV}$.

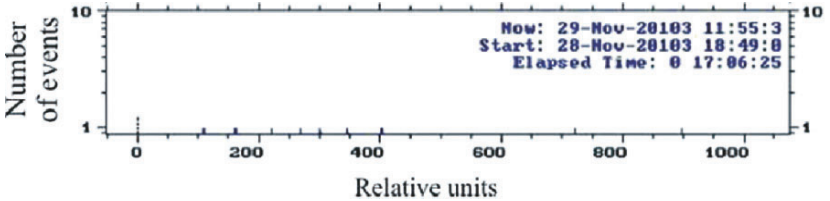


Fig. 10.30. Energy spectrum of α -particles from specimen No. 50 before and after the Cu^{++} action.

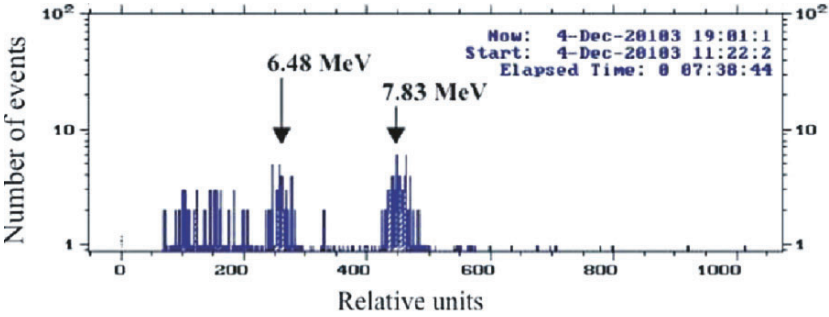


Fig. 10.31. Energy spectrum of α -particles from specimen No. 8383.

Table 10.3. Time dependance data of activity reduction of the induced irradiator.

Spectrum registration mode	Period, hours	Total number of events
Under influence of the ion source	3.0	2572
	7.5	6282
After influence of the ion source	after 2.0	396
	after 10.0	386
	after 17.0	730
	after 18.0	750

The measurement of α -spectra of the specimens irradiated by low-energy Cu ions were carried out in the off-line mode on a VLADO unit with an α -spectrometer positioned in a low-noise lead box constructed and supplied by the “Soyuzenergo” concern. The spectra of background events corresponding to the high-energy α -background are given in Fig. 10.27. The energy resolution of the α -spectrometer was $\Delta E = 16\text{ keV}$ on the lines of ^{226}Ra in the region $E_\alpha = 4\text{ to }8\text{ MeV}$.

It is necessary noting that statistically significant events in the off-line mode were accumulated for 6 to 8 h. Then only the background events caused by high-energy cosmic-ray particles were registered.

The observed α -lines with energies: 6.48 MeV, 7.70 MeV, 7.83 MeV, 8.10 MeV, 8.24 MeV, 9.30 MeV, 10.20 MeV, 10.90 MeV and with the energy resolution $\Delta E = 16\text{ keV}$ by α -lines allow us to assert that the surface of accumulating screens-targets contains the microscopic amounts of SHN which become α -emitters under their low-energy excitation.

In view of the geometry of a detecting system and the period of α -decay of these amounts of SHN, we can estimate their number in micro-clusters.

We estimated the number of SHN in one atomic cluster as 500, and the surface of an accumulating screen contains up to 2000 clusters. The decay α -band (4 to 12 MeV) discovered in one of the experiments allows us to evaluate the parent nucleus mass as 1500 to 2000 a.m.u.

Based on the obtained results, we draw the following conclusions:

1. We have registered superheavy atomic nuclei on the specimens whose surface contains the products of the explosion-induced destruction of a target after a high-energy explosive compression, which is confirmed by their induced fission.
2. The discovered α -decay of superheavy atomic nuclei allows us to evaluate the parent nucleus mass as 1500 to 2000 a.m.u.

For the first time, we have observed α -particles with energies: 6.48, 7.70, 7.83, 8.10, 8.24, 9.30, 10.20, 10.90 MeV as the products of the induced fission of SHN.

10.7. Anomalies of the Spatial Distribution of Extrinsic Elements in the Accumulating Screen and the Synthesis of Superheavy Nuclei

Upon the execution of experiments on the supercompression of a solid target with the help of a coherent electron driver, we discovered the unique spatial distribution of various chemical elements and isotopes with mass numbers in the interval $1 \leq A \leq 240$ in the volume of accumulating screens which are made of a chemically pure element and remote from the collapse zone (see Fig. 11.25–11.27 in Sec. 11.2.6). Most frequently, we used accumulating screens made of Cu (99.99%). It turned out that the synthesized elements and isotopes are positioned, in a number of cases, in the screen volume in spatially coinciding low-dimensional clusters (e.g., Fig. 11.27 in Sec. 11.2.6) lying in the scope of the same extremely thin concentric layer which is symmetrically positioned relative to the screen center at a distance of about $0.3 \mu\text{m}$ from its surface (e.g., Fig. 11.25 in Sec. 11.2.6). At a depth of about $5 \mu\text{m}$, we discovered the other analogous layer. Upon the detailed analysis of these layers, we established that the position of all points of each layer corresponds to the same distance from the screen surface, if we measure it in the direction from HD. Such a position of the layer corresponds to the same braking path of particles arrived from this dot in the screen volume up to the complete stop in the scope of the layer. The situation looks so as if we are faced with the braking of identical heavy particles arrived on the screen

from one point (the focus point corresponding to the place of a collapse in the target volume). It is paradoxical that we registered the completely different elements (light, medium, and heavy ones), rather than hypothetical identical ones!

It is easy to verify that such a distribution over the layer surface and over the distance from the screen surface cannot be a result of the ordinary Coulomb braking of various fast ions fallen on the screen surface from the side of HD. This circumstance is related to that the ions with different masses and ionization multiplicities and with the same initial energy differ very strongly (by several tens of times) in the length of a path to their complete stop in the screen volume.

The simple analysis showed that such a distribution of various chemical elements and isotopes in the scope of the same layer can be realized only if the following conditions are fulfilled:

- All particles braked and stopped in the scope of the layer are identical.
- For the constancy of a charged state of particles, their velocity must be small as compared to the velocity $v_0 = e^2/\hbar = 2.5 \cdot 10^8$ cm/s of valent electrons.
- To ensure a relatively long braking path at a small initial velocity, the mass of an unknown particle M should be very large.
- Various chemical elements and isotopes observed in the scope of the layer are created in the process of nuclear transmutation of these identical and very heavy particles after their stop.

The direct calculation on the basis of a model of elastic collisions was made in Sec. 11.1.2 and showed that the discovered distribution can be derived under various conditions of the experiment only in the case where it is related to identical neutral particles with subatomic size and very large masses $A \approx 1000$ to 8000 . Such particles can be only hypothetical nuclear clusters (superheavy nuclei neutralized by electrons over the volume).

The same analysis of the braking dynamics yields that the initial velocity of motion of these nuclei on the screen surface was comparatively small and equal to $v(0) \approx (2 \text{ to } 4) \cdot 10^6$ cm/s. Such a value ensures the elastic character of the interaction of nuclei with atoms of the target and is simultaneously sufficient in order to ensure a comparatively large path in the target. Different layers correspond to different masses of particles.

The creation of such superheavy nuclei can be connected only with the processes running in the collapse zone in the target volume. After the escape from the collapse zone, they fell on the surface of an accumulating screen and, after the elastic braking, stopped in a concentric layer, in which we registered then various chemical elements. There are weighty reasons to

consider that, after the braking and stop, superheavy nuclei stimulate the running of nuclear reactions with the formation of the whole spectrum of chemical elements and isotopes.

We were able to invent only two scenarios for the running of these reactions.

The first scenario is related to the possible decay of superheavy nuclei after their braking and stop. The correcting experiments showed that such a mechanism was not realized, most likely, in this system. This conclusion is based on the data of experiments with the use of a spectrometer VG-9000, where the total amount of all synthesized extrinsic elements appearing after the single action from the side of the target on the accumulating screen was determined. It turned out that the total mass of all extraneous (synthesized) elements in the screen volume not only exceeds their initial mass in the form of an initial admixture in the screen substance by many orders, but exceeds significantly the mass of all the substance being in the volume of the active zone of the initial target prior to the explosion. We determined the volume of this active zone many times basing on the degree of a decrease in the radioactivity of an analogous target in experiments on the neutralization (liquidation) of radioactivity.

Such an excess of the masses of synthesized elements relative to the total mass of the active zone contradicts the assumption that these extraneous elements are created as a result of the decay of superheavy nuclei formed in the active zone of a target during its explosion.

The second scenario is related to the assumption that the extraneous elements are the secondary products of nuclear transformations of the superheavy nuclei. These transformations can be related to peculiarities of the process of absorption of “ordinary” nuclei of a target by the superheavy nuclei. If such an absorption occurs, then the transmutation leads to an increase in masses of the superheavy nuclei with the help of the nuclear synthesis. The probability of such a synthesis can be sufficiently large due to a high transparency of the Coulomb barrier. This transparency is induced by the neutralization of the charge of a superheavy nucleus by the electrons localized inside it. The process of synthesis is accompanied by the release in energy. During such an absorption of nuclei of a target, the excess of the binding energy can be released by various channels (γ -emission, emission of nuclear fragments, etc.). One of the channels of the energy release can be related to the creation of various “normal” nuclei and to their escape from the volume of a growing superheavy nucleus. It is obvious that such a process is possible if the growth rate for superheavy nuclei will be higher than the creation rate for lighter nuclei. Such a combined “fusion-fission” mechanism is considered in Sec. 11.2.6 and explains practically all the regularities of

the observed phenomenon. Moreover, this mechanism agrees well with results of the above-discussed experiments on the induced decay of superheavy nuclei.

10.8. Substantiation and Discussion of Synthesis and Registration of Superheavy Nuclei

We will try to summarize the above-presented results of several independent cycles of measurements. In our opinion, the totality of these results testifies to the formation of stable (or long-lived quasistable) superheavy nuclei in the setup designed and constructed at the Electrodynamics Laboratory “Proton-21”.

We realize that such a conclusion contradicts many fixed ideas and cannot be accepted without relevant arguments.

It seems to be expedient to clearly separate those principal points which characterize the reliability of the assertion about the reality of the synthesis and registration of superheavy nuclei to the maximum extent and to consider and to discuss those possible counterarguments which can be advanced against the above-presented fundamental assertions.

Counterargument 1: *The results of the presented mass-spectrometric experiments are related to the registration of molecular polyatomic ions, rather than to that of superheavy nuclei with the same total mass.*

By carrying on the spectrometry of superheavy ions and nuclei, we undertook a number of special measures allowing us to prevent the appearance and registration of molecular ions possessing the same total mass, which can lead to false conclusions. These measures are as follows:

- The same specimens were complexly studied on several mass-spectrometers of different types. In particular, the nuclei with mass numbers $277 \leq A \leq 280$ and with $A = 433$ were registered on both a thermoionization mass-spectrometer and a secondary-ion one. The nuclei with charges (Z) 112, 115, 140, 157, and 204 were registered with the help of X-ray spectrum microanalysis and Auger-spectrometry. The nuclei with mass numbers $A \approx 296...310$, 311, 327, 340, 370, and 410 were registered on a mass-spectrometer using the Rutherford backscattering method and devices of other types.
- For a secondary-ion mass-spectrometer SIMS, an analyzer IMS 4f, we used the special “offset” mode of registration which allowed us to separate the molecular complexes and individual monoions with identical mass numbers with a very high efficiency.

- For the same mass-spectrometer, we used a special method of spatial selection of monoions and molecular ions. It consists in using the visualization of a raster of the etching by the primary ion beam and in comparing the probabilities of the simultaneous presence of a hypothetical superheavy nucleus with a specific mass number and the whole collection of those “ordinary” elements, whose interaction can lead to the formation of a molecular ion with the same total mass at the same place of the raster.
- On mass-spectrometers of the other types, the identification of a specific superheavy nucleus was carried out only if the integral spectrum of isotopes contained no ions of those atoms which could form finally a molecule with the same mass as the nucleus of a superheavy ion.
- In a number of control measurements, we used a laser time-of-flight mass-spectrometer, in which the temperature at the laser focus was equal to 100 eV. At such a temperature, the existence of molecular complexes is impossible, and a mass-spectrometer registered only monoions.
- The study of superheavy nuclei was carried out within the method of Rutherford backscattering with the use of accelerated α -particles with an energy of 27.2 MeV and $^{14}\text{N}^{++}$ ions with an energy of 8.7 MeV. Within such a method, only monoions are registered.
- In one of the experiments on a track detector mounted on the pole of a magnet in the vacuum chamber of the magnetic analyzer of collimated beams of the particles emitted by HD of an exploded target, we registered the decay of a pointlike object, in which 276 particles with masses from $A = 4$ to $A = 7$ and energies up to 10 MeV escaped from the one place at the same time.
- We carried out multiple studies of the effect of induced decay on the specimens which contained nonidentifiable superheavy nuclei. In the process of this decay, we registered the formation of superheavy nuclei ($A = 433$), medium ones ($30 \leq A \leq 40$, $A = 63$, $A = 65$), and many light nuclei (mainly α -particles).

The totality of these precautions and their complex character allow us to assert that the registered objects are really stable or quasistable superheavy nuclei.

Counterargument 2: *There are no standard (proton-neutron) nuclear models, basing on which one can compare the results of the performed spectrometric studies of superheavy nuclei with the data of calculations.*

Many such models and calculations are known. As an example, we can mention a number of recent investigations (see Refs. 63, 64, 65) devoted to

the search for the nuclei with masses $A \approx 270$ to 296 in the reactions $\text{Ni}^{64} + \text{Bi}^{209}$, $\text{Zn}^{70} + \text{Pb}^{208}$, $\text{Ca}^{48} + \text{Pu}^{244}$, $\text{Ca}^{48} + \text{Cm}^{248}$ forming the highly excited compound systems with $E^* \approx 50$ to 60 MeV. In the above-mentioned works, it was shown that the formation of the nuclei of SHE in the ground states in the processes with a two-particle input channel is defined by the structure of the shells of colliding atomic nuclei. In this region, the existence of relatively long-lived twice magic nuclei is possible.

The possibility for the existence of superheavy nuclei with significantly greater values of the mass number A and charge Z has already been considered in some works. For example, the work (see Ref. 66) predicts that the magic numbers of neutrons in the region $228 < N < 820$ are as follows: $N = 308$, $N = 406$, $N = 524$, $N = 644$, and $N = 772$.

The corresponding magic numbers of protons, according to that work, are $Z = 210$, $Z = 274$, and $Z = 354$. These results yield that the twice magic nuclei $_{114}\text{X}^{298}$, $_{164}\text{Y}^{472}$, $_{210}\text{Z}^{616}$, $_{274}\text{R}^{798}$ can exist. Upon such magic numbers, the nuclei of SHE must have a higher stability, greater binding energy, and higher abundance as compared to adjacent nuclei. Such nuclei are essentially more stable and closer to the line of β -stability. By comparing these results with those of the study within the method of Rutherford backscattering and by taking into account a large error of the determination of A , we may assert the existence a quite reliable correlation between the predictions of the mentioned work and the results of experiments concerning the region of mass numbers $A \approx 300$, $A \approx 400$, and $A \approx 800$.

Counterargument 3: *The assertion about the synthesis of stable superheavy nuclei contradicts the known data on the decrease (on the average) of the stability of nuclei with increase in masses and charges of nuclei towards the “transuranium region”.*

The main argument, which is usually used as the “proof” of the impossibility for stable superheavy nuclei to exist, is related to the general problem of hydrodynamic stability of heavy nuclei relative to the process of spontaneous fission and to the condition of absolute instability $Z^2/A > 48$ which follows from the analysis of the problem. This condition is derived on the basis of the analysis of the development of a instability of the volume oscillations of a charged liquid drop.

We note that this argument has, in fact, semiphenomenological qualitative character, because it does not take into account a number of peculiarities of the shell structure of a nucleus. The consideration of such a structure leads to the hypothesis of the possibility of the existence of “stability islands” with the simultaneous increase in the lifetime for much more heavier nuclei (in particular, for nuclei with $A \approx 298$). The analysis of some works, which

studied this hypothesis and predicted the possibility of the existence of twice magic nuclei with $A = 472$, $A = 616$, and $A = 798$, is presented above.

At the same time, it is obvious that, according to the traditional proton-neutron model, the stability of nuclei must sharply decrease (on the average) with increase in their masses. This is conditioned by that the forces of repulsion of protons are long-range and sharply increase with the number of protons, whereas the forces of attraction are short-range and weakly depend on the number of nucleons. It is obvious that the processes related to the effect of nucleon shells cannot change radically this situation.

It is worth noting that all the known attempts to solve the problem of creation of superheavy nuclei with the use of the technique of collisions of accelerated heavy nuclei derived on the unique accelerators of heavy particles were quite efficient in the creation of comparatively light transuranium nuclei with $Z < 106$ to 110. As for heavier elements (at $110 < Z < 116$), the efficiency of such a method of synthesis of superheavy nuclei is so low that only single rapidly decaying nuclei were synthesized under conditions of the continuous operation of an accelerator for months. We may conclude that the synthesis of nuclei at the expense of high-energy binary collisions has approached its natural limit. Specific reasons for such a situation are presented in Secs. 11.1.1 and 11.1.2.

The natural question arises: How universal is the idea of the uniqueness of the proton-neutron model of nuclei?

We mention at least two models which can significantly change the idea of the structure of the nuclear matter.

In the first model, the stability of a nucleus is ensured at the expense of the phenomenon of pion condensation. It is named as the “Migdal model” and is thoroughly analyzed in Sec. 11.1.2. This model admits the existence of stable superheavy nuclei with mass numbers up to $A \approx 2 \cdot 10^5$. But such a pion condensate can be realized only under the extremely critical conditions. In particular, it is necessary to preliminarily compress the nuclear matter by 3 to 6 times as compared to the ordinary density of nuclei or to preliminarily create “bare” nuclei with charges $Z \geq \alpha^{-3/2} \approx 1600$. Here, $\alpha = e^2/\hbar c$ is the fine structure constant.

It is obvious that the fulfillment of any of these requirements is the extremely complicated problem under conditions of a terrestrial laboratory.

The second model is much more realistic and is related to the necessity to take into account the stabilizing action of a degenerate relativistic electron gas present in the nucleus volume. In fact, this model rejects the uniqueness of the existence of “pure” superheavy nuclei on the basis of only a proton-neutron system and assumes that a final object can be, in a number of quite real cases, an electron-proton-neutron system, in the scope of

which the complete union of the nucleus and electrons in a single system occurs. In this system, nucleons and electrons are the partners with equal rights and give the quite comparable contributions to the total energy of a nucleus.

The realization of such a system removes at once the majority of problems concerning the existence of superheavy nuclei which seem unsolvable earlier.

In particular, the presence of the degenerate gas of relativistic electrons in the nucleus volume decreases sharply the Coulomb repulsion of protons. In the framework of this quasineutral model, the arguments about the hydrodynamic instability of nuclei lose their sense. In addition, within the model, the interaction of electrons and *nucleons* leads to the self-similar self-compression of an electron-*nucleon* system up to the state of Coulomb collapse and, finally, to a sharp shift of the maximum of the binding energy from the ordinary value $A_{\text{opt}} \approx 60$ to very high values $A_{\text{opt}} \gg 60$ (even up to $A_{\text{opt}} > 3000$). *We believe that just this approach is valid for the the system of many nuclei surrounded by the electron gas, where the Coulomb repulsion of nuclei is strongly suppressed, and the process of collapse can be realized under a certain external action.* This model was intensively discussed in the previous sections.

The additional advantage of this model is related to the absence of the necessity to ensure the particular supercritical preliminary conditions comparable with those occurring in the process of gravitational astrophysical collapse. To start the process of self-compression of the electron-nucleus system of a part of the target, which will lead to the Coulomb collapse, we need only to carry out the action of a certain intensity and a character onto the target which can be easily realized under conditions of ordinary terrestrial laboratories.

Counterargument 4: *It is unclear how the energy gain and comparative simplicity of the synthesis of superheavy nuclei on the used experimental setup can be made consistent with the well-known data, according to which the increase in the mass number of nuclei with a charge more than that of iron is the energy-consuming process.*

This argument is quite valid if we restrict ourselves to only such nuclei which are described by the “standard” proton-neutron model. For such a system, the synthesis of heavy nuclei requires the large expenditures of energy. This energy is mainly spent for the execution of the work on overcoming the Coulomb repulsion, i.e., on bringing together many protons in the nucleus volume.

The situation will change if we consider other types of the nuclear matter.

In particular, the proton-neutron-pion model of nuclei (the Migdal model) is characterized by the dependence of the mean potential energy of each nucleon on the mass number which reveals two potential wells, i.e., two minima: the first minimum is “ordinary” with $A_{\text{opt}} \approx 60$ in the region near iron, and the second one is “anomalous” with $A_{\text{opt}} \approx 200\,000$. These potential wells are separated by a very high and wide potential barrier, whose position corresponds to the charge $Z \approx 1600$. These potential wells are symbolically presented in Sec. 11.1.1 in Fig. 11.1.

Since the bottom of the second well is deeper by 15 to 20 MeV than that of the first one, any nucleus “having passed” from the region of “ordinary” nuclei to the region of “anomalous” superheavy nuclei upon the formation of the pion condensate turns out unstable relative to the increase in A and can absorb “ordinary” nuclei many times by increasing its mass up to A_{max} . It is natural that this process is energy-gained and leads to the release of a great energy.

For this system, the main problem is related to the overcoming of the potential barrier separating two regions. We have shown earlier that if the certain starting conditions for the process of collapse are satisfied, the very strong and correlated Coulomb interaction of a degenerate electron gas and nuclei will lead to the deformation, lowering, and disappearance of the potential barrier separating two regions and to a shift of the position of the first minimum from the initial value $A_{\text{opt}} \approx 60$ to the side of very large values $A_{\text{opt}} \gg 60$. Of great importance is the circumstance that these starting conditions can be by many orders weaker by their energy contribution than the energy which will be released during the self-similar self-controlled process of self-compression of a target, which leads to the final complex of nuclear transformations.

Counterargument 5: *If the synthesis of superheavy nuclei is energy-gained, how can you explain the phenomenon of a stimulated decay of these nuclei which is energy-consuming in the framework of the considered model?*

These processes are completely consistent. One of the simple explanations of such a consistency consists in that the increase in the mass of a superheavy nucleus occurs at the expense of the absorption of adjacent “ordinary” nuclei by it. Since this process will be energy-gained, it is accompanied by the release of a great binding energy of the electron-nucleus system and by the heating of the nucleus. Different ways of the energy release can be realized. One of the most efficient ways to cool the superheavy nucleus “starting to boil” is the cluster decay with the formation of an additional

comparatively light nucleus on the surface of the growing superheavy nucleus. In the framework of such a process, the absorption of one or several nuclei of a target with the total mass number A_Σ leads to the simultaneous creation (emission) of one or several nuclei with the total mass number A_F . Such a process of the “fusion-fission” type is possible if $A_F < A_\Sigma$. In this case, the corresponding conservation laws are satisfied.

Counterargument 6: *It is well known from the practice of nuclear physics that, for the running of nuclear reactions with the participation of nuclei with large charges, a great kinetic energy of relative motion of these nuclei is necessary. This contradicts the conditions for the realization of the above-considered experiments on the induced decay which require a very small energy (on the scale of nuclear physics).*

The analysis of the numerous experiments on the stimulation of decays shows that we observed the phenomena which are unambiguously related to the processes of nuclear transformations. This is confirmed by a high efficiency of the use of several independent and purely nuclear methods (in particular, the method of track detectors and detectors of α -particles). These transformations occur only on specimens (accumulating screens) subjected to the direct action of the explosion products dispersed from HD. The very small (on the scale of the processes characteristic of nuclear physics) energy of charged particles (probing electrons with an energy of 30 keV and ions with an energy of 10 to 12 keV) or infrared optical quanta acting on the target and heating its ions at the irradiated place up to a temperature of several or several tens of eV allows us to unambiguously conclude that such an action cannot directly induce the running of electronuclear, photonuclear, or barrier-involved nuclear reactions accompanied by the creation of α -particles and heavier ions with a great energy.

Still to the greater extent, these arguments concern the processes leading to the formation of thin layers in the accumulating screen volume with different concentrations of different elements. In this case, the kinetic energy of the relative motion of hypothetical superheavy nuclei and atoms of a target was extremely small and did not exceed 1 eV/nucleon.

Both the theory of nuclear reactions and the practice of applied nuclear spectroscopy do not admit the running of the nuclear reactions of fission on a large scale upon the action of charged particles with such small energy on nuclei.

In our opinion, these processes can run in the nuclei of SHE. It is obvious that these nuclei (or, at least, a part of them) are in a metastable state. We may assume that at least two alternative mechanisms of such processes are realized.

On the one hand, the reactions of decay can run in the case where these nuclei are far from the stability line. In this case, we expect the existence of a mechanism ensuring a single decay of nuclei upon a relatively weak external action (e.g., at the expense of the removal of the spin exclusion or other one). The proper estimates of the results of experiments show that such a mechanism of decay is improbable. This follows from the fact that, upon the evolution of nuclei being far from the stability line, most probable are β -processes which should lead to the emission of electrons and positrons and to the generation of the accompanying characteristic emission, rather than the processes of emission of nuclear clusters. Near the region where such reactions stimulated by the external action are running, we must observe a very strong background of X-rays and γ -emission. However, neither essential γ -emission nor considerable β -activity were found in the experiments on stimulated decay.

On the other hand, similar reactions can run by the following reason. Though these superheavy nuclei are on the stability line (which can be different from the stability line of "ordinary" nuclei), but their binding energy is far from the maximum value. In this case, we must assume that a weak external action on the nearest environment of a nucleus of SHE will stimulate the reactions involving the interaction of a nucleus of SHE with the adjacent nuclei of the accumulating screen surface. Due to a very small thickness of the Coulomb barrier near the nuclei of SHE, the probability of such reactions will be large. These reactions can lead to the stage-by-stage synthesis of the heavier nuclei of SHE and to the accompanying emission of α -particles and other registered nuclei.

Because the one-type reactions with the creation of α -particles and heavier nuclei run upon completely different actions (at the expense of the interaction with electrons, ions, and optical quanta), it is obvious that the ways of stimulation of the possible mechanisms of these reactions will not be purely nuclear (possibly, that they are not nuclear at all!). In fact, a weak external action in this case must fulfill the function of a distinctive catalyst. The generalizing analysis of all experimentally realized methods of stimulation of the induced decay shows that all they are combined by only one total characteristic, namely the intense local heating of a small part of the specimen under study. This circumstance allows us to assume that the local thermal perturbation and the increase in the mean kinetic energy of the nuclei of a target play a very important role in the stimulation of nuclear reactions with participation of the nuclei of SHE. We have developed a physical model of thermal stimulation of the nuclear reactions involving the superheavy nuclei and nuclei of a target. Now it passes the experimental verification.

Thus, the above-considered justified objections and counterarguments, which are advanced against the fundamental conclusion on the reality of the discovery of stable superheavy nuclei at the Electrodynamics Laboratory “Proton-21”, are naturally answered, agree sufficiently well with the developed theoretical models, and do not disprove the very conclusion.

The authors realize that their conclusions and arguments will not be uniquely perceived by all readers and potential critics, but think that the arguments by totality have already surpassed the critical mass which is required for the sure assertion as for the reality of the existence of a controlled artificial synthesis of superheavy nuclei and their discovery.

The authors understand also that the theoretical models, which substantiate the possibility of the existence of stable superheavy nuclei, will be obligatorily reconsidered and supplemented in the future. Fortunately, the process of cognition is infinite.

Part IV

Preliminary Résumé of Obtained Results,
Theories, and Physical Models

STABILITY OF ELECTRON-NUCLEUS FORM OF MATTER AND METHODS OF CONTROLLED COLLAPSE

S. V. Adamenko and V. I. Vysotskii

11.1. Controlled Electron-Nucleus Collapse of Matter and Synthesis of Superheavy Nuclei

11.1.1. General Problems of Synthesis of Superheavy Nuclei

The study of extremum states of matter is one of the most important trends of contemporary physics. Prior to the construction of superhigh-energy accelerators, this region of physics was referred to astrophysics. Only in astrophysical objects, we would expect the presence of such conditions, under which the fundamental changes in the structure and properties of matter can occur. Such states of matter are not restricted to a certain spatial interval and can be realized on the subatomic and subnuclear scales, as well as in the limits of macroscopic regions. It seemed to be almost obvious that only the action of the gravitational forces leading to the gravitational collapse of massive stars can ensure these conditions. The situation is significantly changed after the construction of accelerators producing elementary particles, ions, and atomic nuclei with extremely high energies and, recently, of lasers with extremely small duration of pulses of induced emission and, respectively, extremely high intensity of such pulses. The parameters of such drivers allow one to hope for the achievement, with their help, of such pressure and temperature that exist in the volume of white dwarfs and in the core of a neutron star due to the gravitational driver. At the same time, it seems to be obvious that any artificial drivers cannot ensure such fundamental transformations of matter which occur, for example, in the core of a neutron star.

If we digress from so fundamental comparisons, it becomes evident that, under the terrestrial conditions, the problem of a controlled transformation of the nuclear matter and, in particular, the problem of creation of anomalous and superheavy nuclei are of great importance. These problems have scientific and important applied meaning. They are key in the solution of such applied problems as the reactor-free usage of atomic (nuclear)

energy, technological synthesis of rare isotopes, utilization of the worked-out nuclear fuel, and development of precise analog methods of the nuclear technology on the basis of nuclear chemistry.

For the sake of justice, it is necessary to note that, up to now, there is no sufficiently convincing solution of the problems of the synthesis of superheavy nuclei in astrophysical objects.

The attempts to solve this problem under conditions of a terrestrial laboratory cannot be recognized successful as well. The problem of the creation of superheavy nuclei with the use of the technique based on collisions of accelerated heavy nuclei derived on the unique accelerators of heavy particles is investigated in several laboratories of the world and requires huge financial and material-technical expenditures. Though the method of collisions was sufficiently efficient upon the creation of relatively light transuranium nuclei with $A < 110$, it became very inefficient in the attempts to create nuclei with $A > 116$. We can conclude that the method of synthesis of nuclei at the expense of collisions has approached its natural limit. The extremely low efficiency of this method of creation of superheavy nuclei and the related possibility of fatal errors in the interpretation of results (e.g., Refs. 34, 35) are based on the basically unavoidable flaws of the method. It is quite obvious that the method of synthesis of superheavy nuclei in counter-collisions is unsuitable for the technological purposes and, all the more, upon the solution of the above-mentioned problems.

Alternative is the mechanism of formation of superheavy nuclei upon the use of a stimulating action of the pion condensate on the stability of the nuclear matter (Refs. 36–38).

A brief analysis of the peculiarities and efficiency of the collision-based method of creation of superheavy nuclei and the method of pion condensation will be given in Sec. 11.1.2.

In the subsequent sections and in the following chapter, we consider a basically other scenario for the creation of superheavy nuclei that does not require collisions of nuclei and does not use the phenomenon of pion condensation. This scenario is based on our works and is founded on the phenomenon of stimulation of the self-controlled process of formation of a collapse of the electron-nucleus system. Such a process is adapted to the electron-nucleus structure of matter. It leads to the displacement of the maximum of the binding energy of this system to the region of large mass numbers and can ensure both the barrier-free synthesis of superheavy nuclei and the accompanied synthesis of the extremely wide spectrum of “ordinary” nuclei, whose considerable part possesses a significantly distorted isotope ratio (sometimes by tens and hundreds of times as compared to the natural ratio of the same isotopes).

We have first shown that there exists the threshold value of the electron density of a compressed atom such that its achievement induces the self-controlled process of “falling of electrons onto the nucleus” and the formation of an extremely compressed electron-nucleus cluster. The lowest threshold corresponds to nuclei with high charge. These questions will be considered in Sec. 11.1.3.

The considered mechanism of radical self-controlled reconstruction of an electron-nucleus system allows one to answer several principal questions:

1. How can the energy gain of the synthesis of superheavy nuclei be achieved?
2. What does the stability of superheavy nuclei depend on?
3. Which way is the high efficiency of the synthesis of such nuclei achieved in?
4. How can one agree the starting low energy and power of the used driver with such a spectrum of fundamental transformations of the electron-nucleus matter that cannot be derived upon the use of other, incomparably more intense laboratory drivers?

In addition, this mechanism allows us to satisfy those critical requirements (in particular, those of the synthesis of nuclei with very large charges Z), whose fulfillment can lead to the realization of the phenomenon of pion condensation in nuclei. The analysis of these processes on the basis of the mechanism proposed by us is carried out in Sec. 11.1.4.

In Sec. 11.1.5, we consider the peculiarities of a manifestation of the collapse effect of an electron-nucleus plasma in astrophysical objects. The starting conditions for such processes can be realized relatively easily in the process of the gravitational collapse of a massive star and, under certain assumptions, in the volume of white dwarfs. In addition, these processes can be directly related to the nucleosynthesis of various nuclei, including neutron-deficient nuclei, and to such global effects as the burst of a supernova and the shaking of its shell.

In Sec. 11.1.6, we consider the threshold conditions and peculiarities of the running of such base reactions as the generation of γ -emission accompanying the process of compression, creation of electron-positron pairs, neutronization and protonization of nuclei in the zone of a collapse at the stage of its formation.

In our opinion, the considered mechanism corresponds sufficiently well to the process of synthesis of stable superheavy nuclei that are observed since 2000 at the Electrodynamics Laboratory “Proton-21” in the experiments performed without the use of nuclear reactors and accelerators of heavy ions.

11.1.2. Problems and Prospects of the Creation of Superheavy Nuclei from Heavy Particles Collisions and with the Help of Pion Condensations

Problems of the Creation of Superheavy Nuclei in Collisions of Heavy Particles. Below, we give a brief analysis of the collision-based method of synthesis of superheavy nuclei. It is necessary to note that the main conception lying in the base of the method of creation of superheavy nuclei with the use of colliding heavy particles bears a frankly “force” character. It is related to the assumption that, in a collision of two heavy nuclei with a high energy of relative motion, it is possible to overcome the Coulomb barrier and to form a compound nucleus with the total mass close to that of colliding nuclei. Despite the high scientific-technical level of the available accelerating equipment, such a method reminds ideologically the procedure of derivation of fire upon the impact of two stones that was used by primitive men. It is easy to prove that the method of creation of superheavy nuclei based on this conception is extremely inefficient and limited from the conceptual viewpoint.

We can give several arguments to confirm the above conclusions.

1. The very low probability of the central collision of nuclei accelerated to a high energy to overcome the Coulomb barrier leads to a very low probability of the formation of a compound nucleus. In view of the ratio of the cross sections of synthesis of superheavy nuclei ($\sigma_f \approx 10^{-10}$ b) and that of ionization of atoms of a target ($\sigma_i \approx 10^8$ b), it is obvious that the dominant part of the kinetic energy of such nuclei is expended for the ionization of atoms of a target.
2. The heavy compound nucleus formed in the process of synthesis from two lighter nuclei with their proton-neutron ratios close to the “normal” one, $1/(\alpha - 1) \approx 0.8$ to 0.6 , turns out inevitably to be neutron-deficient and has a great excess of protons. For this reason, it turns out unstable due to the significant Coulomb repulsion of protons. Here, $\alpha = A/Z \approx (2.2 \text{ to } 2.5)$ is the ratio of the mass and charge number of nuclei on the line of stability.
3. The very large energy of relative motion of the colliding nuclei, that played a positive role and was quite necessary on the stage of overcoming the Coulomb barrier, plays a negative role and turns out excessive in the formed compound nucleus. The presence of this energy induces the overheating of the nucleon system of the compound nucleus, fast emission of nucleons and clusters, and the decay of the nucleus.
4. The very small cross section of the synthesis make the use of counterbeams to be extremely inefficient and leads inevitably to the

necessity to use condensed dense immovable targets irradiated with fast nuclei. For such a geometry of experiments, the necessity to satisfy the momentum conservation law leads to that the synthesized superheavy nuclei are received a very high kinetic energy (hundreds of MeV) and interact in the flight mode with the rest nuclei of a target. Such an interaction causes the additional instability of the compound nuclei and a lot of secondary background effects complicating the study and identification of these nuclei.

These circumstances lead finally to a very small number of formed and registered superheavy nuclei. The typical rate of registration of the nuclei with $Z \approx 112$ to 118 by the products of their decay does not exceed one nucleus for one month (against the general background of events in the system of registration, being at least 10^{11} to 10^{12} for the same period). It is natural that, for such an efficiency of registration, any detailed study of such nuclei and, all the more, their usage are basically impossible. This leads, in particular, to very great errors upon the determination of the lifetime of such nuclei.

Moreover, these factors become more and more essential with increase in the charge and mass of synthesized superheavy nuclei. Indeed, for the synthesis of heavier synthesized nuclei with the final charge Z and atomic number $A = \alpha Z$, it is necessary to use heavier starting nuclei with the total charge $(Z_i + Z_j) \geq Z$. To overcome the Coulomb barrier upon the interaction of such nuclei, a greater kinetic energy of relative motion, which is proportional to the product of the charges of colliding nuclei, $T \sim Z_i Z_j$, is required. After the formation of a compound nucleus, the mean energy of excitation per one nucleon,

$$\langle T_1 \rangle \approx T/\alpha(Z_i + Z_j) \sim Z_i Z_j/\alpha(Z_i + Z_j), \quad (11.1)$$

grows proportionally to the ratio $Z_i Z_j/(Z_i + Z_j)$.

For example, for the synthesis of a nucleus with $Z = 118$ by the assumed reaction $\text{Pb}^{208} + \text{Kr}^{86} \rightarrow \text{X}^{294}$, the irradiation of a solid target made of Pb^{208} was carried out on a cyclotron at Berkeley by the beam of Kr^{86} ions with an energy of $E_{\text{Kr}} \approx 450 \text{ MeV}$ or $\Delta E_{\text{Kr}} \approx 5.2 \text{ MeV/nucleon}$ (Refs. 34, 35). Upon the interaction of these two nuclei, the laws of conservation of energy and momentum yield that, in the case of their full fusion, the final kinetic energy of the compound nucleus is equal to a lesser part ($\eta \equiv M_{\text{Kr}}/(M_{\text{Pb}} + M_{\text{Kr}}) \approx 29\%$) of the initial energy E_{Kr} of the incident particle. The rest of energy, $(1 - \eta)E_{\text{Kr}} \approx 318 \text{ MeV}$, is spent on inelastic intranuclear processes (an increase in the energy of the ground state of a compound nucleus and its heating).

As follows from the drop nuclear model, an increase in the energy of the ground state of a compound nucleus (it is numerically equal to a decrease in the binding energy of nucleons) relative to the sum of the energies of starting nuclei is

$$\begin{aligned}
 E_{\text{exit}} = & \varepsilon_2(A_{294}^{2/3} - A_{\text{Kr}}^{2/3} - A_{\text{Pb}}^{2/3}) \\
 & + \varepsilon_3 \left[(Z^2/A^{1/3})_{294} - (Z^2/A^{1/3})_{\text{Kr}} - (Z^2/A^{1/3})_{\text{Pb}} \right] \\
 & + \varepsilon_4 \left\{ [(A/2 - Z)^2/A]_{294} - [(A/2 - Z)^2/A]_{\text{Kr}} - [(A/2 - Z)^2/A]_{\text{Pb}} \right\} \\
 & + \varepsilon_5 (1/A_{294}^{3/4} - 1/A_{\text{Kr}}^{3/4} - 1/A_{\text{Pb}}^{3/4}) \approx 276 \text{ MeV},
 \end{aligned} \tag{11.2}$$

which is about 61% of the initial energy E_{Kr} .

Here, $\varepsilon_2 = 17.8 \text{ MeV}$, $\varepsilon_3 = 0.7 \text{ MeV}$, $\varepsilon_4 = 94.8 \text{ MeV}$, and $\varepsilon_5 = 33.57 \text{ MeV}$ are the parameters of the drop nuclear model for the considered interaction of even–even nuclei.

The remaining part of the kinetic energy of an incident particle (about 10% or 45 MeV) defines the strong heating and excitation of the compound nucleus, which led to its instantaneous decay in most cases.

In very rare cases, the compound nucleus can remain unbroken by releasing the excess energy in the process of emission of several neutrons. Numerous estimations show that the probability of the “survival” of an overheated compound nucleus is at most 10^{-10} . At the same time, it is obvious that because the initial compound nucleus is neutron-deficient at the moment of its formation, such a version of the “survival” with the emission of neutrons is far from the optimum one since it additionally increases this deficit. These circumstances make the synthesis of superheavy nuclei with the optimized proton–neutron ratio on the basis of the collision-based technology to be impossible.

It is clear that these circumstances impose a fundamental limitation on the synthesis of superheavy nuclei and make the process of formation of even though individual nuclei with $A \geq 300$ on the basis of such a collision-based mechanism to be improbable. It is also obvious that such a way of the synthesis is basically unsuitable for the derivation of a macroscopic amount of matter on the basis of superheavy nuclei and, all the more, for its use in technology.

It is worth noting that the mentioned sources of the instability of synthesized superheavy nuclei are inherent, in the first turn, in the very method of their creation and have no relation to the problem of stability of superheavy nuclei with the optimized composition. For this reason, the characteristics of few successful experiments on the synthesis of superheavy nuclei (in fact, the synthesis of individual nuclei on the unique accelerators)

cannot be adequately compared with those of the known models of such nuclei. Such a situation hampers the development of the theory of nuclear matter and leads to a certain canonization of the nuclear structure as being only in the form of pure proton-neutron ensembles not admitting the presence of other particles in it.

It is proper to note that the method of synthesis of superheavy nuclei with $Z > 100$ on the basis of the reactor-based technology (upon the irradiation of heavy stable nuclei by neutrons) is also extremely inefficient due to the availability of many alternative ways for the transmutation and fast decay of activated nuclei in the interval of the acts of the subsequent capture of neutrons. Even under astrophysical conditions, a similar mechanism of synthesis of superheavy nuclei at the expense of the fast subsequent capture of a large number of neutrons (the r -process) does not lead to the synthesis of superheavy nuclei.

It is obvious that these circumstances are among the reasons, due to which the questions of the creation and stability of superheavy nuclei with $A > 250$ to 300 on the basis of alternative models of the nuclear structure turned out to be weakly studied.

An alternative approach to this problem was considered in the cycle of works of A. Migdal (Refs. 36–38) devoted to the problem of the creation of superheavy nuclei with the use of the phenomenon of condensation of pions in the volume of a nucleus. The use of the pion condensate allows one, in perspective, to shift the maximum of the binding energy of a nucleus to the region of very large mass numbers $A \gg 60$, which opens the way to the formation of superheavy nuclei without external action. At the same time, for the formation of such a condensate, we need the conditions that are considered unattainable at present.

Prospects, problems, and threshold conditions of the synthesis of superheavy nuclei with the use of the phenomenon of condensation of pions in the volume of a nucleus. In their works carried out in the 1970–1980s, A. Migdal and his collaborators (Refs. 36–38) considered the premises of the creation of quasistable superheavy nuclei with charges and mass numbers in the intervals from $Z \geq (\hbar c/e^2)^{3/2} \approx 1\,600$ and $A \geq 3\,000$ to $A \geq 200\,000$ and of anomalous superdense nuclei with $A \approx 2Z \geq 100$. As the main mechanism of the stability of such nuclei, Migdal considered the influence of the condensate of pions in the volume of a nucleus.

In the simple form, the idea of the influence of pions on the stability of nuclei looks as follows. Pions stabilizing superheavy nuclei can be formed in the reaction of spontaneous decay $n \rightarrow p + \pi^-$. This reaction can run only in a very deep potential well formed by protons in two types of nuclei: in superdense nuclei with the density of nucleons exceeding the “normal”

one by 3 to 6 times or in nuclei with a large charge $Z \geq (\hbar c/e^2)^{3/2} \approx 1600$. Migdal called the latter type of nuclei “supercharged”. The difference of the Fermi energies of the proton and neutron subsystems of a nucleus, that arises upon the decay of neutrons, is eliminated at the expense of the β^+ decay of protons, $p \rightarrow n + \beta^+ + \nu$. Positrons can leave the nucleus and annihilate upon the interaction with electrons of the target.

The accumulation of negative pions leads to their condensation on the lower energy level. The obvious influence of pions on the nucleus stability is realized, in fact, through two mutually consistent mechanisms.

First of all, negative pions screen a part of the Coulomb potential created by protons. However, it is obvious that because the total number of protons exceeds the number of pions, such a mechanism cannot ensure the total screening and the final stability of a nucleus. The final stability of a nucleus in the presence of the pion condensate is ensured by the specific quantum-mechanical square-law effect of attraction of protons that is realized by pions. This attraction is independent of the charge sign and occurs in the case where the energy of a particle (a pion) in the very strong Coulomb field of protons in a nucleus exceeds its rest energy $m_\pi c^2$. At the ordinary mean distance between protons in a nucleus $\langle r_{pp'} \rangle \approx 1.3 \times 10^{-13} (A/Z)^{1/3}$ cm, such an attraction can be ensured only by sufficiently light particles (in particular, pions). It is obvious that the attraction of a pion to several adjacent protons is equivalent to the mutual attraction of these protons. On the contrary, heavier protons being on the same distance cannot ensure such a nonlinear mode of the mutual attraction (for them, the threshold of a manifestation of the nonlinear interaction is revealed at the distance considerably smaller than the distance between adjacent protons). Hence, every pion stabilizes twice a superheavy nucleus: first, as a negatively charged particle, and, secondly, as a lighter (as compared to a proton) particle participating in the specific quantum-mechanical attraction.

Moreover, because the final influence of the screening of protons by negative pions is weaker than the influence of the quantum-mechanical square-law effect of attraction that does not depend on the charge sign of pions, the stability of a superheavy nucleus can be ensured not only by negative pions, but by positive pions as well.

The additional complication of these processes, which does not change the main result (the stabilization of nuclei), is related to the collective character of the interaction of protons with the Bose-condensate of pions.

It is worth noting that the efficiency of the considered quantum-mechanical attraction of nucleons increases with decrease in the mean distance between nucleons in a nucleus or with increase in the total charge of a nucleus. In view of the scenario under consideration, it is obvious that

the efficiency of the action of the pion-based mechanism of stabilization of nuclei increases upon the additional compression of a nucleus (i.e., upon the increase in both the density of nucleons and the Coulomb field in a nucleus) or for a very large initial charge of a nucleus. In this case, the increase in both the intranuclear Coulomb field and the Coulomb repulsion of protons is compensated with excess by the stronger effect of quantum-mechanical attraction with the participation of pions. The analysis performed by Migdal yields that the energy gain of the process of self-compression of a nucleus can lead to a final increase in the density of nucleons in a nucleus by a factor of 3 to 6 up to the value corresponding to the balance of the considered forces of attraction and repulsion of nucleons at very small distances. This repulsion is resulted from a strong correlation of the positions of nucleons and is usually considered in the frame of the model of independent pairs.

The presence of the stabilizing influence of the pion condensate allows one to predict the stability of nuclei with pion condensate in the region of the masses of nuclei up to $A_{\text{opt}} \geq 200\,000$. Moreover, because the question about the existence of a finite limit of A_{max} in the frame of the Migdal model is not studied up to now, it can turn out that such a limit does not exist!

The general structure of the dependence of the nucleus energy on the mass number with regard the existence of the pion condensate (per one nucleon and by normalizing on a Fe nucleus) is presented in Fig. 11.1. The first potential well with minimum at $A_{\text{opt}} \approx 60$ corresponds to region 1 of “ordinary” relatively light nuclei with $A < 250$. This region is separated

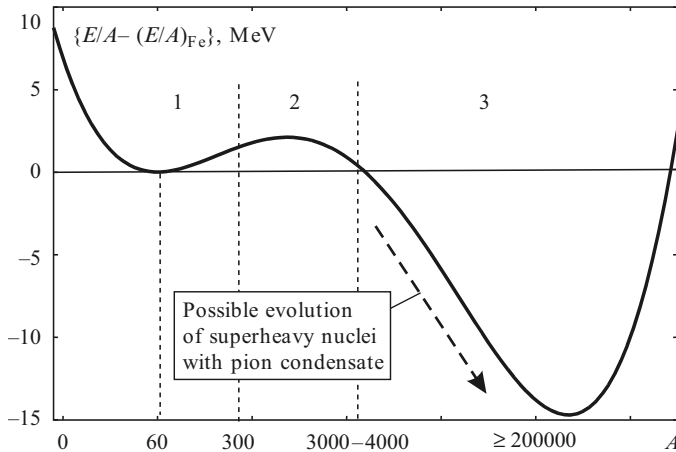


Fig. 11.1. Conjectural scheme of the dependence of the nucleus energy on the mass number A at the normal nuclear density and in the absence of a pion condensate in the nucleus volume (region 1) or in its presence (region 3).

by a very high barrier (region 2) from the second, considerably wider and deeper well with minimum at $A_{\text{opt}} \geq 200\,000$ (region 3). Since the bottom of the second well is by 15 to 20 MeV deeper than that of the first one, any nucleus “passed” from region 1 in region 3 turns out to be unstable with increase in A upon the formation of the pion condensate and can multiply absorb “ordinary” nuclei, by increasing its mass up to A_{max} .

The main question consists in that how to overcome the energy barrier in region 2 and to reach the state of pion condensate. Migdal believed that the height of this barrier can turn out to be so large that the Universe existence time is insufficient for a spontaneous tunnel transition to be realized. A certain external action is needed. It follows from the Migdal model that the transition across region 2 from the first to the second well is possible on the basis of any of the two conditions: in the presence of a very large charge ($Z \geq 1600$) in the initial nucleus or under the initial pulse uniform compression of a nucleus up to the state corresponding to an increase in the density of nucleons by 3 to 6 times. If we assume that the same coupling law for A and Z is true in the region of the second potential well as in the region $A < 300$, the mentioned charges must correspond to nuclei with $A \approx 3400$ to 4000 .

Obviously, it is very difficult to satisfy these two conditions.

We note at once that it is meaningless to discuss the first condition of the formation of a pion condensate on the basis of the hypothesis on the possibility to create the initial nuclei with $Z \geq 1600$ by using the collision-based method of synthesis (especially in the light of the above-mentioned flaws of such a synthesis).

One of the possible ways to satisfy the second condition, i.e., to reach the initial compression of the nuclear matter by 3 to 6 times for the realization of the pion condensation was considered by B. Pontecorvo (see the reference in Ref. 36). He proposed to use the process of collision of heavy accelerated nuclei for the compression of a nucleus (e.g., nuclei of Pb).

It is easy to calculate the energy needed for such a compression and independent of a specific mechanism.

In the degenerate nonrelativistic gas of nucleons in a nucleus, the Fermi energy of nucleons $E_{F0} = (3\pi)^{2/3}(\hbar^2/2m_n)n_n^{2/3} \approx 40$ MeV and their mean energy $\langle E_{10} \rangle = (3/5)E_{F0} \approx 24$ MeV are connected with the density of nucleons n_n by the relation $\langle E_1 \rangle \sim E_F \sim n_n^{2/3}$. These results yield that, in order to increase n_n by 3 to 6 times, it is necessary to increase the mean energy of nucleons up to $\langle E_1 \rangle \approx (3 \text{ to } 6)^{2/3} \langle E_{10} \rangle \approx 50$ to 90 MeV/nucleon in the process of uniform compression of the nuclear matter. This value should be increased by $k \approx 2$ to 3 times with regard to the fact that only a part of the kinetic energy of relative motion is transformed into the energy of

nucleons upon the collision of nuclei. In this case, the Fermi energy increases up to $E_F \approx 84$ to 132 MeV. In the very optimistic case, we would expect that the shock waves arising upon the collision can induce the necessary enhancement of the nuclear matter density at the center of a formed compound nucleus with a short lifetime.

This scenario needs several serious corrections.

The enhancement of the density up to the required value in the process of action of a shock wave can be reached only in the case of the 3D axisymmetric compression of the compound nucleus. It is obvious that such a condition cannot be fulfilled upon the pairwise collision of two nuclei due to the absence of a converging shock spherical wave in this case. The necessary (but not sufficient) condition for the excitation of such a wave is the presence of an axisymmetric synchronized external action on a nucleus. The mean energy of such an action (per one nucleon) is $\Delta E_1 = (\langle E_1 \rangle - \langle E_{10} \rangle) > (26 \text{ to } 66)k \text{ MeV/nucleon} \approx (50 \text{ to } 200) \text{ MeV/nucleon}$.

Respectively, the total energy of the external action (per one nucleus) is $\Delta E = A\Delta E_1 \approx (5 \text{ to } 11)k \text{ GeV/nucleus} \approx (10 \text{ to } 33) \text{ GeV/nucleus}$ (at $A \approx 200$).

With regard to the fact that the mean velocity of nucleons in a compressed nucleus,

$$\langle v \rangle = (3/5)^{1/2} v_F = (6E_{F0}/10m_n)^{1/2} \approx 10^{10} \text{ cm/s}, \quad (11.3)$$

differs slightly from the light velocity, the process of compression of a nucleus with radius R_n must run at most $\Delta t \approx R_n/\langle v \rangle \approx 10^{-22} \text{ s}$. The same duration must correspond to the considered external action on a nucleus. If the external action is longer, then the nucleons accelerated at the leading edge of the external action will leave the nucleus prior to the time when the common potential well confining all the nucleons of the compressed nucleus is formed. This condition yields that the minimum power spent on the compression of every nucleus corresponds to an unreal great value $\Delta P \geq \Delta E/\Delta t \approx 10^{13} \text{ W/nucleus}$ (at $A \approx 200$). Such a power corresponds to an improbably great intensity of an external action, $\Delta J \geq \Delta P/4\pi R_n^2 \approx 10^{37} \text{ W/cm}^2$.

It is obvious that a similar energy contribution in the form of axisymmetric action on a nucleus is impossible (at least, on the modern level of technology). For the sake of comparison, we indicate that the best sources of coherent action (powerful femtosecond lasers) can provide the intensity at a level of at most $\Delta J \approx 10^{20} \text{ W/c}^2$.

In addition, it is obvious that this effect (even if such a value of ΔJ is realized on the nucleus surface) will not lead, most likely, a compressed cold nucleus as the required result. The point is that the given estimations

assume that all the energy of an external driver is fully spent on the increase in the Fermi energy of the degenerate gas of nucleons without increase in its temperature. We note that, in the process of compression at the expense of the use of a self-focusing (upon the axisymmetric action) shock wave in the medium of a degenerate Fermi-gas, the main effect is an increase in the medium temperature (Ref. 39) rather than an increase in the density. In this case, the greater value of ΔJ is necessary. In addition, the heating of nucleons related to a shock wave leads at once to their emission from a nucleus by the scenario identical to that of the collision-based synthesis.

Thus, at the expense of a hypothetical force action upon the collision of nuclei, it is impossible to satisfy the conditions for the supercompression of a nucleus and for the formation of a pion condensate in nuclei of the ordinary type with small charge.

The presented estimations yield that though the pion condensation on the basis of “supercharged” nuclei according to the primary idea does not require a collision of fast nuclei and can run without expenditures of energy, but the very necessity to use the initial nuclei with charges $Z \geq 1600$ for the formation of such a condensate allows us to consider this model as a purely hypothetical one.

Hence, the method of pion condensation in its primary form is unable to ensure the synthesis of superheavy nuclei with $100 < Z < 1600$ and to realize the stage-by-stage transition of stable or long half-life “normal” nuclei from region 1 characteristic of ordinary nuclei with $Z \leq 100$ across the high potential barrier 2 to region 3 with $Z \geq 1600$, where the existence of nuclei with pion condensate is possible. At the same time, such a method can be very efficient upon the synthesis of nuclei heavier than those with $Z \approx 1600$.

The question about which mechanism can ensure both the synthesis of superheavy nuclei in the interval $100 < Z < 1600$ and the “transition” from region 1 to region 3 will be analyzed in what follows.

11.1.3. Mechanism and Threshold Conditions for Heavy Nuclei Formation in Degenerate Electron Plasma

Coulomb Interaction of a Nucleus with Electrons as a Basic Mechanism of the Stimulation of the Synthesis of Superheavy Nuclei. It is obvious that mechanism of stimulation of the reaction of synthesis must explain, in a consistent way, the possibility of the process of synthesis of superheavy nuclei, its high rate, and the stability of synthesized nuclei relative to the decay. It is obvious that the optimum mechanism of stimulation of the synthesis of superheavy nuclei must satisfy the threshold conditions in the region of charges $100 < Z < 1600$ by means of the stimulation of the processes of internal self-organization of matter on the

nuclear and subnuclear levels, rather than by a force action on a nucleus that would be commensurate with the expected total effect. In this case, such a situation seems to be ideal, when the role of the technological factors of the external action is revealed only during the first phase (the phase of a “bare” stimulation). Then the process of global reconstruction of the nuclear matter would develop by the proper laws of self-similar systems with the use of inner energy sources. Within such an approach, the role of the first phase is analogous to that of the first photon which begins the process of laser generation.

In our opinion, in the definite range of conditions, the electron subsystem plays the defining role in the process of inner self-organization of matter on the nuclear and subnuclear levels, which corresponds to a “force-free” method of solution of the problem of synthesis of superheavy nuclei and defines a greater probability of the synthesis.

This assertion is based on the following several facts:

1. The electron system of every nucleus is maximally adapted to this nucleus due to the conditions of adiabaticity, which allows us to consider the electron-nucleus interaction to be ideally symmetric, consistent, and efficient.
2. The laws of quantum mechanics yield that electrons in the field of a nucleus being in the “atomic form of matter” are always in the state of inversion and possess a very large potential energy relative to the nucleus even on the ground energy level of any atom.
3. Under certain conditions, this energy can be released in the process of “falling of electrons onto the nucleus” and can be used for the internal self-organization of matter under its transition from the “atomic form of matter” to the state of collapse of the degenerate electron-nucleus plasma both in the volume of every atom and in an ensemble of atoms.

Below, we consider a “force-free” mechanism of the electron stimulation of the process of self-controlled nuclear transformations which has high efficiency in the region $100 < Z < 1600$ by overcoming the barrier separating the “ordinary” nuclei from the nuclei with a spontaneously formed pion condensate. In addition, the mechanism under consideration can efficiently act at lesser values Z (including $Z \leq 92$).

Consider the peculiarities of this mechanism in more details.

The problem of evolution of any nuclear systems is usually considered on the basis of models taking into account the strong interaction of the nearest nucleons in the volume of a nucleus and the Coulomb interaction of all protons. In nuclear physics, the directivity of a reaction (in the direction of synthesis or fission) is considered in view of its energy gain.

The nucleus energy E_n can be expressed with sufficient accuracy and generality (without exact definition of the structure of nucleon shells in the frame of the drop nuclear model) with the help of the Weizsäcker formula containing the sum of terms, each of which corresponds to one kind of interaction,

$$E_n = Zm_p c^2 + (A - Z)m_n c^2 - \varepsilon_1 A + \varepsilon_2 A^{2/3} + \varepsilon_3 Z^2/A^{1/3} + \varepsilon_4 (A/2 - Z)^2/A + \varepsilon_5/A^{3/4}. \quad (11.4)$$

Here, the quantity $\varepsilon_1 A$ is the binding energy of all nucleons; $\varepsilon_2 A^{2/3}$ is the surface energy of a nucleus, $\varepsilon_3 Z^2/A^{1/3}$ is the Coulomb energy of repulsion of protons in a nucleus; $\varepsilon_4 (A/2 - Z)^2/A$ is the symmetry-related energy taking into account the increase in the energy upon breaking the equality of the numbers of protons and neutrons because of the necessity to populate higher energy levels due to the satisfaction of the Pauli principle separately for neutrons and protons; and $\varepsilon_5/A^{3/4}$ is the term taking into account the pairing effects for even–even, odd–odd, and odd–even nuclei.

Many correcting versions of this formula are known, but it is characteristic that, in all the cases, the possibility for an essential influence (not in the form of a small perturbation) of the external electron environment on nuclear processes was not taken into account. As a rule, such a viewpoint is conditioned by that the binding energy of electrons with a nucleus is slight on the scale of intranuclear processes if the density of electrons is small.

Quantum mechanics explains the reason for such a situation.

Consider two extreme cases with different degrees of the influence of a nucleus on the motion of electrons.

In the system including nuclei and electrons not subjected to an external action, only those quantum-mechanical states are possible which correspond to the solution of the equation of Schrödinger or Dirac involving the Coulomb field of the nucleus. These equations yield that the maximum binding energy of an electron with the nucleus,

$$E_{\max} = Z^2 e^4 m_e / 2\hbar^2, \quad (11.5)$$

corresponds to the deepest level from the allowed ones (level 1s in atoms). For the heaviest stable nuclei ($Z \approx 100$), the binding energy of 1s electrons does not exceed $E_{\max} \approx 200$ keV. The binding energy of the remaining electrons is significantly lower (by 4 and more times). In atoms of the ordinary kind (at $Z < \hbar c/e^2 \approx 137$), the binding energy of electrons with a nucleus cannot exceed value (Eq. 11.5). It is obvious that, in such a system, the influence of electrons on the balance of forces defining the evolution of a nucleus is slight. It is also obvious that such a situation is a direct consequence of the

initial statement of the problem: one considers the system consisting of a nucleus interacting with free electrons by the Coulomb law (it is the principal interaction). In this case, the influence of the environment (the external pressure on the electron shell) is considered to be weak and insignificant as compared to the Coulomb interaction in the atom.

Consider the basically other, alternative case. Let a system of nuclei and electrons exist, on which the so strong external pressure acts that namely it defines the behavior of electrons. In this situation, the interaction of electrons with nuclei is weaker and introduces small corrections to results of the influence of the external pressure.

The transition between the considered cases of weak and extremely strong external actions corresponds to several sequential events.

At the zero or very weak external pressure, electrons and nuclei exist in the form of bound atoms. Upon an increase in the pressure, there occurs the destruction of the outer electron shells of atoms. The outer shells are completely destroyed at such a density of atoms, when the distance $2R$ between their nuclei becomes equal to the doubled radius $r_1 = \hbar^2/m_e e^2 \approx 0.5 \rho A$ of outer electron shells. This corresponds to the ion density $n_i \approx (m_e e^2/2\hbar^2)^3 \approx 10^{24} \text{ cm}^{-3}$. Upon the further increase in the pressure and the attainment of the ion density of the compressed plasma $n_{icr} \approx 10^{25} Z \text{ cm}^{-3}$, the majority of inner atomic shells is broken (see Ref. 40). At a higher pressure, to which the ion density $n_i \geq 10Z(m_e e^2/\hbar^2)^3 \approx 10^{26} Z \text{ cm}^{-3}$ corresponds, there occurs the full destruction of the atomic structure for the heaviest atoms on all levels including the deepest ones. This state corresponds to a nonrelativistic electron-ion plasma with electron density $n_e \geq 10Z^2(m_e e^2/\hbar^2)^3 \approx 10^{28}$ to 10^{30} cm^{-3} .

This plasma is degenerate by the electron component, if its temperature is less than the degeneration temperature $T_{\text{deg}} \approx 350 \text{ keV}$ equal to the Fermi energy $E_F = (3\pi^2)^{2/3}(\hbar^2/2m_e)n_e^{2/3}$. At a higher compression (at the density $n_e \geq 10^{31} \text{ cm}^{-3}$), the electron component of the plasma will correspond to the relativistic degenerate electron gas and, at $n_e \gg 10^{31} \text{ cm}^{-3}$, to the superrelativistic gas. We note that the density $n_e \approx 10^{30}$ to 10^{31} cm^{-3} , being very large on the scale of the existence of ordinary atoms (of which the value $\langle n_e \rangle \approx 10^{23}$ to 10^{24} cm^{-3} is characteristic), is simultaneously very small as compared to the density of nucleons $n_n \approx 10^{38} \text{ cm}^{-3}$ in a nucleus.

The neutral electron-nucleus plasma derived as a result of the external compression of some volume of matter is degenerate by the electron component and is characterized by several processes of internal interaction of its components:

1. In a nucleus, there occurs the interaction of nucleons. The total energy of this interaction is defined by five last terms in Eq. 11.4 for E_n .

2. Nuclei interact with their close electron environment and attract electrons to themselves. In addition, nuclei interact one with another, which favors the creation of a long-range order in the electron-nucleus plasma.
3. Electrons interact one with another in correspondence with their spin and charge.

Consider two last kinds of interaction in more details.

At the high electron density, the minimum energy of such a system, with regard to the interaction of the charges of nuclei, corresponds to the formation of a bcc lattice of nuclei in the degenerate electron plasma.

If the density of electrons neutralized by nuclei corresponds to the condition $n_e \gg Z^2(m_e e^2 / \hbar^2)^3 \approx 10^{25} Z^2 \text{ cm}^{-3}$, then the behavior of electrons corresponds to the ideal degenerate Fermi-gas (see Ref. 41).

Consider the electron-nucleus plasma as a system consisting of periodically arranged neutral Wigner-Seitz cells with a nucleus at the center of every cell (Fig. 11.2). Every of the cells can be approximately described as a sphere. The radius of every cell, $R = (3Z/4\pi n_e)^{1/3}$, is determined from the condition of its electric neutrality (the number of electrons in it is equal to the number of protons Z in a nucleus located at the center of a cell).

The boundary of every cell undergoes the action of the pressure compressing the plasma and keeping it from a dispersion. This pressure is a result of the summary action of different mechanisms, the main of which are the Coulomb attraction of electrons to the nucleus, mutual Coulomb

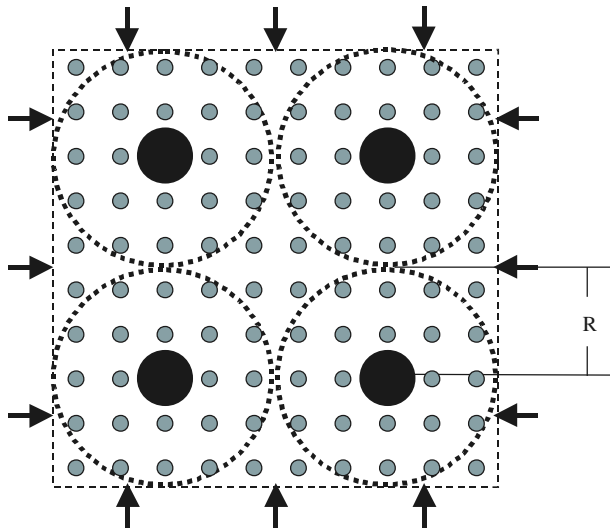


Fig. 11.2. System of nuclei with charge Ze in the degenerate electron gas with density n_e .

repulsion of electrons, and kinetic (Fermi) pressure of the degenerate electron gas. Every Wigner–Seitz cell is a distinctive compressed quasiautom, in which the state of electrons is mainly formed at the expense of the action of this pressure. In the first approximation, we can neglect the electrical interaction of different cells. We note that the account of such an interaction at a very high density leads inevitably to the formation of a cubic lattice of nuclei surrounded by the degenerate relativistic gas of electrons.

Can the interaction of all particles in the volume of a neutral elementary Wigner–Seitz cell with the degenerate electron gas cause the stability of a compressed quasiautom or its self-compression in the absence of the external pressure? Such a question was first considered in some works (e.g., Refs. 42, 43) upon solving the problem of the stability of such space objects as black holes, white dwarfs, and neutron stars).

In a number of works (e.g., Refs. 42–45), it was shown that the presence of the electron-nucleus attraction in the preliminarily compressed degenerate nonrelativistic electron gas leads to a decrease in the Fermi pressure of the electron gas by several percents, which is clearly insufficient for the stabilization of the plasma and, all the more, for its self-compression. For the sake of justice, it is worth noting that the estimations of the influence of the Coulomb interaction of a compressed relativistic degenerate gas were also made in Ref. 43, which led to the same negative answer. However, these estimations cannot be recognized as reliable, because the formula for the Coulomb interaction energy $U = -e\varphi(r) = -Ze^2/r$ that corresponds to the nonrelativistic case was used in the calculations on the basis of the equation of state of a relativistic degenerate gas. Below, this question will be considered in more details.

On the other side, a detailed analysis of the interaction of a nucleus with the relativistic degenerate gas of electrons without the presence of an additional external boundary (i.e., without a preliminary external compression of a Wigner–Seitz cell) was performed in Refs. 37, 38. The author of those works concluded that, in a nucleus with $Z \ll 1600$, the presence of electron shells corresponds to an atom with a very great density of electrons in the volume of the nucleus. Beyond the scope of the nucleus, this density (and the corresponding wave function of electrons) asymptotically decreases to zero upon the unlimited increase in the distance. In fact, these boundary conditions define the final structure of a free atom with very large charge Z .

It is obvious that such a calculation corresponds to a single free atom with large charge Z , but not to the very large periodic system of strongly compressed Wigner–Seitz cells (i.e., a periodic system of quasiautoms) bounded in volume. The surface of contact of these cells ensures the

corresponding boundary conditions which are characterized by a nonzero wave function of electrons and the zero derivative of this function.

Thus, the question about the evolution of a quasiatom preliminarily compressed by the external forces up to the state of a relativistic degenerate gas with a sharp boundary of separate quasiatoms turns out to be insufficiently studied. The analysis performed below will show that just such a system can pass, under certain conditions, to the state of self-controlled electron-nucleus collapse, whose evolution promotes the formation of super-heavy nuclei.

Interaction Energy of the Compressed Relativistic Degenerate Gas of Electrons and a Nucleus in the Volume of a Neutral Wigner–Seitz Cell. Now we calculate the interaction energy of all particles around a nucleus in the volume of a neutral elementary Wigner–Seitz cell.

First, we calculate the interaction energy of one electron with the nucleus and other electrons.

In the general case, the motion of an electron in the field of a potential $\varphi(r)$ is described by the Dirac equation

$$\{E + e\varphi(r) - \hat{\beta}m_e c^2 - \hat{\alpha}c\mathbf{p}\}u = 0, \quad (11.6)$$

$$\hat{\beta} \equiv \hat{\alpha}_4 = \begin{pmatrix} \hat{1} & 0 \\ 0 & -\hat{1} \end{pmatrix}, \quad \hat{\alpha} = \begin{pmatrix} 0 & \hat{\sigma} \\ \hat{\sigma} & 0 \end{pmatrix}, \quad (11.7)$$

for the four-component wave function u .

Here, $\hat{\beta}$ and $\hat{\alpha}$ are the scalar and vector Dirac matrices defined through the 4-component unit matrix $\hat{1}$ and the 4-component vector Pauli matrix $\hat{\sigma}$.

Let us multiply the equation by the operator $\{E + e\varphi(r) - \hat{\beta}m_e c^2 - \hat{\alpha}c\mathbf{p}\}$ and use the operator identities

$$(\hat{\sigma}\mathbf{A})(\hat{\sigma}\mathbf{B}) = (\mathbf{A}\mathbf{B}) + i\hat{\sigma}[\mathbf{A}\mathbf{B}], \quad (\hat{\sigma}\mathbf{A})(\hat{\sigma}\mathbf{A}) = \mathbf{A}^2 \quad (11.8)$$

and also the permutation relations for components of the vector Dirac matrix

$$\hat{\alpha}_i \hat{\alpha}_k + \hat{\alpha}_k \hat{\alpha}_i = 2\hat{\delta}_{ik}. \quad (11.9)$$

Then we get a differential second-order equation of the Schrödinger equation type:

$$\{\hat{\mathbf{p}}^2/2m_e + U_{\text{eff}}\}u = \varepsilon u, \quad (11.10)$$

$$\varepsilon = [E^2 - (m_e c^2)^2]/2m_e c^2, \quad (11.11)$$

$$U_{\text{eff}} \equiv U_{\text{eff}}(r, E) = -e\varphi(r)(E/m_e c^2) - [e\varphi(r)]^2/2m_e c^2. \quad (11.12)$$

Here, $U_{\text{eff}}(r, E)$ is the effective potential energy defining the motion of an electron with total energy E in the field of a potential $\varphi(r)$.

Eq. 11.12 for the effective potential energy $U_{\text{eff}}(r, E)$ differs essentially from the “traditional” expression for interaction energy, $U_0(r) = -e\varphi(r)$ valid only in the nonrelativistic case (at $E \approx m_e c^2$ and $(E - m_e c^2)/m_e c^2 \ll 1$) and in relatively weak fields ($|e\varphi(r)| \ll m_e c^2$).

In particular, for nonrelativistic particles with mass m and full energy $E \approx mc^2$ ($(E - mc^2)/mc^2 \ll 1$), the effective potential energy in a very strong field is determined by the first two terms in Eq. 11.12 and looks like

$$U_{\text{eff}} \approx U_{\text{eff}}(r, mc^2) \approx -e\varphi(r) - [e\varphi(r)]^2/2mc^2. \quad (11.13)$$

Equation 11.13 says that, for nonrelativistic charged particles in the case of a strong scalar potential (at $|e\varphi(r)| > 2mc^2$), the negative contribution of the second (“nonlinear”) term exceeds always the action of the first (“linear”) term. The last can have different signs depending on the charge sign. Such a result leads to the final “nonlinear” attraction ($U_{\text{eff}}(r, mc^2) < 0$) independent of the charge sign of a particle. This effect was discussed above upon the analysis of peculiarities of the interaction of condensed pions with protons in a nucleus in the Migdal model of pion condensate.

In the case of relativistic charged particles, besides a nonlinear attraction, we are faced with the increase in the effective scalar potential connected with the Lorentz transformation of the field of a scalar potential $\varphi(r)$ for rapidly moving particles.

With regard to the relativistic formula for the total energy of one electron

$$\begin{aligned} E &= -e\varphi(r) + (p^2 c^2 + m_e^2 c^4)^{1/2} \\ &= -e\varphi(r) + T_1(p) + m_e c^2 \\ &\equiv -e\varphi(r) + \gamma_p m_e c^2, \end{aligned} \quad (11.14)$$

Eq. 11.12 yields the formulas for the effective potential energy $U_{\text{eff}}(r, p)$ of a specific electron with momentum p and for the analogous effective energy $U_{\text{eff}}(r, p_F)$ of an electron on the Fermi-sphere surface (with the limiting momentum p_F):

$$U_{\text{eff}}(r, p) = -e\varphi(r) - e\varphi(r)[T_1(p) - e\varphi(r)/2]/m_e c^2 = -e\varphi(r)\gamma_p + [e\varphi(r)]^2/2m_e c^2, \quad (11.15)$$

$$U_{\text{eff}}(r, p_F) = -e\varphi(r)E_F/m_e c^2 + [e\varphi(r)]^2/2m_e c^2 = -e\varphi(r)\gamma_F + [e\varphi(r)]^2/2m_e c^2. \quad (11.16)$$

Here, $E_F = (p_F^2 c^2 + m_e^2 c^4)^{1/2} = \gamma_F m_e c^2$ is the Fermi energy, $T_1(p) = (p^2 c^2 + m_e^2 c^4)^{1/2} - m_e c^2 = (\gamma_p - 1)m_e c^2$ is the kinetic energy of an electron

with momentum p in the composition of the degenerate relativistic electron gas, γ_p and γ_F are, respectively, the relativistic Lorentz-factor for an electron with momentum p and with maximum momentum p_F on the Fermi-sphere surface.

We note the basic difference of the formula for the effective potential energy of an electron in the quasiatome compressed by the external pressure from that for an ordinary atom kept by only the internal Coulomb interaction.

In bound electron-nucleus systems without external action (without additional external pressure), the total energy $E = -e\varphi(r) + T_1(p) + m_e c^2$ is always less than the rest energy $m_e c^2$. This follows obviously from the fact that a stable state of an electron bound in an atom is impossible at $-e\varphi(r) + T_1(p) > 0$. In view of the condition $E < m_e c^2$, Eq. 11.12 yields that the maximum (by absolute value) of the potential energy of an electron in the ordinary atom is always bounded:

$$|U_{\text{eff(atom)}}(r, p_F)|_{(\text{max})} \leq e\varphi(r) + [e\varphi(r)]^2/2m_e c^2. \quad (11.17)$$

In the case of a compressed quasiatome, the total energy of a single electron E depends on the external pressure which keeps the compressed quasiatome from a dispersion and can unlimitedly increase its kinetic energy $T_1(p) = (\gamma_p - 1)m_e c^2$, which corresponds to the unlimited increase in the ratio $E/m_e c^2$. We note that, upon compression, it is possible to increase the kinetic energy up to values $\langle T_1(p) \rangle \gg \langle e\varphi(r) \rangle$ which are basically unattainable in the Coulomb field, for which we have $\langle T_1(p) \rangle = \langle e\varphi(r) \rangle/2$ according to the virial theorem.

Such a compression results in a sharp increase in the effective potential energy of an electron in the compressed quasiatome:

$$|U_{\text{eff(quasiatome)}}(r, p)|_{(\text{max})} > e\varphi(r) + [e\varphi(r)]^2/2m_e c^2. \quad (11.18)$$

In particular, if $E \gg \langle e\varphi(r) \rangle$, then

$$|U_{\text{eff(quasiatome)}}(r, p)|_{(\text{max})} \gg e\varphi(r) + [e\varphi(r)]^2/2m_e c^2, \quad (11.19)$$

which corresponds to a very significant increase in the interaction of an electron with the nucleus. The physical reason for the increase in the interaction is related to a relativistic transformation of the scalar potential for moving particles.

Just this circumstance allows us to realize the below-considered effect of self-compression of the electron-nucleus plasma which can occur under a preliminary threshold compression of matter. In the case of "ordinary" atoms not compressed by external forces up to the state of degenerate relativistic gas, such an effect is basically impossible.

We note one more important circumstance.

In the presence of other electrons in the volume of a Wigner–Seitz cell, the potential $\varphi(r)$ [along with the effective potential U_{eff} following from it] must be defined by the potential of the nucleus $\varphi_0(r) = Ze/r$ and by the mean potential formed by other electrons. The summary concentration of electrons $n_e(r) = (1/3\pi^2\hbar^3c^3)[E_F - U_{\text{eff}}(r, p_F)]^3$ in the volume of a Wigner–Seitz cell is a result of such a self-consistent interaction.

The problem of the determination of such a self-consistent potential $\varphi(r)$ can be solved with the use of the relativistic Thomas–Fermi equation

$$\Delta\varphi(r) = -4\pi e \left[[E_F - U_{\text{eff}}(r, p_F)]^3 / 3\pi^2\hbar^3c^3 - Z\delta(\mathbf{r}) \right] \quad (11.20)$$

supplemented by specific boundary conditions.

As mentioned above, such a problem (for the boundary condition $r\varphi(r) \rightarrow 0, r \rightarrow \infty$ and for the simplified potential [Eq. 11.13]) was considered by Migdal (see Ref. 36) upon the analysis of the characteristics of the system of electrons in the field of a nucleus with an arbitrary charge Ze . It is obvious that such a boundary condition corresponds to an isolated free atom uncompressed by external forces, on whose boundary (as $r \rightarrow \infty$) the potential energy $-e\varphi(r)$, density electrons $n_e(r)$, and electron pressure P_e are zero. This model leads to the presence of a screened Coulomb potential $\varphi(r) = \chi(r)Ze/r$ in the volume of a “Thomas–Fermi atom”.

The function $\chi(r)$ satisfies the boundary conditions $\chi(r) \rightarrow 1$ as $r \rightarrow 0$ and $\chi(r) \rightarrow 0$ as $r \rightarrow \infty$ and looks

$$\chi(r) = \exp(-r/R_{S(0)}), \quad R_{S(0)} \approx 0.885(\hbar^2/m_e e^2)/Z^{1/3}. \quad (11.21)$$

In such a “natural” isolated atom, whose size and structure are defined by only the interaction of the nucleus and electrons, the effect of self-compression is impossible.

We will consider the basically other model, namely the system of compressed quasiatoms, in which the size and structure of every atom are defined on the initial stage by the external pressure, and the surface of their contact satisfies the corresponding boundary conditions characterized by a nonzero wave function of electrons and the zero derivative of this function.

It is also necessary to note that some of the parameters of an extremely compressed degenerate electron gas in the presence of a separate (isolated) nucleus in its volume were studied in a number of works earlier. The screening length $R_S = 5.8(3/4\pi n_e)^{1/3}$ of a separate nucleus in the infinite-volume degenerate relativistic electron gas was calculated in Refs. 46, 47 with the use of the static dielectric function $\epsilon_e(k, 0)$ found in Ref. 48. This result is inapplicable for the below-considered compressed quasiatom with

an “ordinary” nucleus and can be used only in the case where the screening length R_S is much less than the quasiautom radius $R = Z^{1/3}(3/4\pi n_e)^{1/3}$ (i.e., at $Z \gg 200$).

The model of a compressed atom of finite size on the basis of the Thomas–Fermi equation was considered in Ref. 44. This model yields, in particular, that the influence of the screening effect becomes more and more insignificant with increase in both the degree of compression and the concentration of electrons and nuclei.

It was noted Refs. 42–44 that, in a compressed medium with mass density $\rho \gg 10^4 \text{ g/cm}^3$ (this corresponds to the electron density $n_e \gg 3 \times 10^{27} \text{ cm}^{-3}$), the state of a degenerate electron plasma corresponds to the ideal homogeneous gas with density n_e independent of a coordinate, and the influence of the Coulomb interaction can be taken into account as the interaction of this homogeneous gas with the field of the nucleus. The same result directly follows from Eq. 11.20 while decreasing R .

These circumstances significantly simplify the further calculations.

The total Coulomb potential energy $U_Q = U_{QL} + U_{QNL}$ of all Z degenerate electrons in the volume $V = 4\pi R^3/3$ of the considered compressed quasiautom with radius R and the nucleus with charge Ze consists of the linear [relative to $e\varphi(r)$] part, U_{QL} , and the nonlinear one, U_{QNL} [proportional to $(e\varphi(r))^2$], and can be derived from Eq. 11.16 by summing the contributions of all electrons with regard to their distribution in the scope of the Fermi-sphere with the limiting (Fermi) momentum $p_F = (3\pi^2)^{1/3} \hbar n_e^{1/3}$.

The linear part of the potential energy is equal to

$$\begin{aligned} U_{QL} &= \sum_{i=1}^Z U_{\text{eff}}(r_i, p_i) \\ &= \int_V \int_0^{p_F} U_{\text{eff}}(r, p) (d^2 N / dp dV) dp dV = W_Q K_F. \end{aligned} \quad (11.22)$$

Here,

$$W_Q = - \int_V e\varphi(r) n_e dV \quad (11.23)$$

is the total Coulomb energy of electrons with regard to their homogeneous distribution in the volume of the compressed quasiautom, but without regard for the influence of their motion in the Fermi condensate;

$$K_F \equiv \langle \gamma \rangle = \int_0^{p_F} (p^2 c^2 + m_e^2 c^4)^{1/2} p^2 dp / \pi^2 \hbar^3 n_e m_e c^2 \quad (11.24)$$

is the coefficient of dynamic growth of this energy related to the motion of electrons; and

$$d^2N/dpdV = p^2/\pi^2\hbar^3 \quad (11.25)$$

is the density of the states of electrons in the r -space and p -space in the degenerate electron gas normed on the total number of electrons

$$\int_V \int_0^{p_F} (d^2N/dpdV) dp dV = Z. \quad (11.26)$$

The coefficient K_F is equal to the mean value of the Lorentz-factor $\langle\gamma\rangle$ for all electrons of the compressed quasiatome. For the relativistic degenerate gas of electrons, $\langle\gamma\rangle = (3/4)\gamma_F$.

The first part W_Q of the total energy U_{QL} in Eq. 11.22 consists of the sum of two components: the energy of the mutual attraction of the nucleus and electrons:

$$\begin{aligned} W_{Q(en)} &= \int_0^R V_0(r) n_e 4\pi r^2 dr = -(3/2)(Z^2 e^2/R) \\ &= -(3/2)(4\pi/3)^{1/3} Z^{5/3} e^2 n_e^{1/3}, \end{aligned} \quad (11.27)$$

and the energy of the mutual repulsion of electrons:

$$\begin{aligned} W_{Q(ee)} &= \int_0^R Q(r) dQ(r)/r = (3/5)Z^2 e^2/R \\ &= (3/5)(4\pi/3)^{1/3} Z^{5/3} e^2 n_e^{1/3}. \end{aligned} \quad (11.28)$$

Here, $Q(r) = -Zer^3/R^3$ is the total charge of all homogeneously distributed degenerate electrons located in a sphere with radius $r < R$ in the volume of the compressed quasiatome.

The coefficient of dynamic growth of the total Coulomb energy is

$$\begin{aligned} K_F &= \left(m_e^3 c^3 / 8\pi^2 \hbar^3 n_e \right) \{ (p_F/m_e c) [2(p_F/m_e c)^2 + 1] \\ &\quad \times [(p_F/m_e c)^2 + 1]^{1/2} - \text{Arsh}(p_F/m_e c) \}. \end{aligned} \quad (11.29)$$

For the nonrelativistic degenerate gas (at $p_F \ll m_e c$ and $n_e < 10^{30} \text{ cm}^{-3}$),

$$K_F = 1 + (3/10)(3\pi^2)^{2/3} (\hbar/m_e c)^2 n_e^{2/3} \approx 1 + 0.3(n_e/10^{30} \text{ cm}^{-3})^{2/3}. \quad (11.30)$$

For the superrelativistic gas (at $p_F \gg m_e c$ and $n_e \gg 10^{30} \text{ cm}^{-3}$),

$$K_F = (3/4)(3\pi^2)^{1/3}(\hbar/m_e c)n_e^{1/3} \approx 0.77(n_e/10^{30} \text{ cm}^{-3})^{1/3}. \quad (11.31)$$

Respectively, the nonlinear part of the total Coulomb potential energy reads

$$\begin{aligned} U_{QNL} &= \int_V [e\varphi(r)]^2 n_e dV / 2m_e c^2 \\ &= \int_0^R (-Ze^2/r + Ze^2 r^2/R^3)^2 n_e 4\pi r^2 dr / 2m_e c^2 \\ &= [8\pi Z^{7/3} e^4 (3/4\pi)^{1/3} / 7m_e c^2] n_e^{2/3}. \end{aligned} \quad (11.32)$$

A change in the total Coulomb energy of the degenerate gas of electrons Eqs. 11.27 and 11.28 is accompanied by a simultaneous change in the exchange energy of the degenerate electron gas (see Ref. 49)

$$\begin{aligned} U_{ee,\text{exch}} &= -(e^2/n_e) \int_0^{p_F} d^3 p_1 / (2\pi\hbar)^3 \times \int_0^{p_F} d^3 p_2 / (2\pi\hbar)^3 4\pi\hbar^2 / |\mathbf{p}_1 - \mathbf{p}_2|^2 \\ &= -[3(3\pi^2)^{1/3} / 4\pi] Z e^2 n_e^{1/3} \end{aligned} \quad (11.33)$$

and by increase in the kinetic (Fermi) energy

$$\begin{aligned} U_{eF} &= \int_V \int_0^{p_F} T_1(p) (dN^2/dp dV) dp dV \\ &= (V/\pi^2 \hbar^3) \int_0^{p_F} [(p^2 c^2 + m_e^2 c^4)^{1/2} - m_e c^2] p^2 dp \\ &= Z m_e c^2 (K_F - 1). \end{aligned} \quad (11.34)$$

The total energy of the degenerate electron gas is

$$\begin{aligned} U_{e\Sigma} &= U_{QL} + U_{QNL} + U_{ee,\text{exch}} + U_{eF} \\ &= [W_{Q(en)} + W_{Q(ee)} + Z m_e c^2] K_F + U_{ee,\text{exch}} - Z m_e c^2 + U_{QNL} \\ &= \{Z m_e c^2 - (9/10)(4\pi/3)^{1/3} Z^{5/3} e^2 n_e^{1/3}\} K_F \\ &\quad - [3(3\pi^2)^{1/3} / 4\pi] Z e^2 n_e^{1/3} - Z m_e c^2 \\ &\quad + [8\pi Z^{7/3} e^4 (3/4\pi)^{1/3} / 7m_e c^2] n_e^{2/3}. \end{aligned} \quad (11.35)$$

Based on the derived results, it is easy to prove that, for the relativistic degenerate gas, the nonlinear component of the energy U_{QNL} for nuclei

with $Z \ll 2800$ turns out to be significantly less than the linear component U_{eQL} and can be omitted for such nuclei.

Evolution and the Process of Collapse of a Neutral Wigner–Seitz Cell at a Great Density of the Compressed Relativistic Degenerate Gas of Electrons. Equation 11.35 yields the effective total pressure

$$P_e(Z, n_e) = -dU_{e\Sigma}/dV = (n_e^2/Z)dU_{e\Sigma}/dn_e. \quad (11.36)$$

of the electron gas on the surface of a compressed quasiatom.

In the case of the ultrarelativistic gas of degenerate electrons and upon the use of the approximate relation $K_F \approx (3/4)(3\pi^2)^{1/3}(\hbar/m_e c)n_e^{1/3}$, the pressure of the electron gas takes the form:

$$P_e(Z, n_e) \approx (3\pi^2/64)^{1/3} \hbar c n_e^{4/3} - [9\pi(4^{1/3}/20)(\hbar^2/m_e c)Z^{2/3} - (16\pi/3)(3/4\pi)^{1/3}(e^4/7m_e c^2)Z^{4/3}]n_e^{5/3}. \quad (11.37)$$

In this formula, the first (positive) term is proportional to $n_e^{4/3}$ and corresponds to the kinetic Fermi pressure of the degenerate gas directed from the center of a Wigner–Seitz cell.

The second term depends stronger on the density of the degenerate electron gas, is proportional to $n_e^{5/3}$, and characterizes the pressure induced by the Coulomb interaction. The direct analysis of this term shows that the negative Coulomb pressure of the degenerate superrelativistic electron gas on the surface of a Wigner–Seitz cell in the direction to its center (i.e., to the nucleus) can be attained for any nuclei with

$$Z < Z_{\max(Q)} = [189(16\pi/3)^{1/3}/320](\hbar c/e^2)^{3/2} \approx 2400 \quad (11.38)$$

(including the “ordinary” nuclei with $Z < 100$).

As a result, we get that the negative total pressure $P_e(Z, n_e) < 0$ of the degenerate superrelativistic electron gas can be realized for any nuclei with $Z < Z_{\max(Q)}$ if the threshold density of the electron gas n_e satisfies the main condition

$$n_e > \left\{ (3\pi^2/64)^{1/3} \hbar c / [9\pi(4^{1/3}/20)(\hbar^2/m_e c)Z^{2/3} - (16\pi/3)(3/4\pi)^{1/3}(e^4/7m_e c^2)Z^{4/3}] \right\}^3. \quad (11.39)$$

In particular, at $Z \ll Z_{\max(Q)}$, this condition becomes

$$n_e > (3000/23328\pi)(m_e c^2/e^2)^3/Z^2 \approx 3 \times 10^{36} Z^{-2} \text{ cm}^{-3}. \quad (11.40)$$

It is expedient to note that the derived maximum value $Z_{\max(Q)} \approx 2400$ defines the upper limit of the charge of a nucleus in the super-compressed medium based only on the influence of the electron Fermi-condensate. Because the value $Z_{\max(Q)}$ exceeds the minimum threshold value $Z_{\min(\pi)} \approx (\hbar c/e^2)^{3/2} \approx 1600$ which is required to start the process of pion condensation in the volume of a nucleus in the Migdal model, the possible further growth of the charge of a nucleus Z beyond the scope of the region $Z < Z_{\max(Q)}$ will be realized already on the basis of the model of pion condensate.

For the more correct calculation of the threshold of a Coulomb collapse on the basis of the conditions

$$P_e(Z, n_e) = -dU_{e\Sigma}/dV = 0, \quad (11.41)$$

it is necessary to use the exact formula for K_F (Eq. 11.29) in the whole range of variations in the electron gas density, and not just in the limiting cases of nonrelativistic and superrelativistic degenerate gases. Below, we carry out the numerical analysis of this problem.

The dependence of the energy $U_{e\Sigma}$ (Eq. 11.35) of the degenerate electron gas on its density n_e for nuclei with $Z = 92, 82, 70, 50, 26, 20, 16$, and 14 per one electron is presented in Figs. 11.3, 11.4, and 11.5 for two subsequent ranges of variations in the density of the degenerate electron gas: $n_e = 10^{24}$ to 10^{29} cm^{-3} and $n_e = 10^{31}$ to $3 \times 10^{34} \text{ cm}^{-3}$.

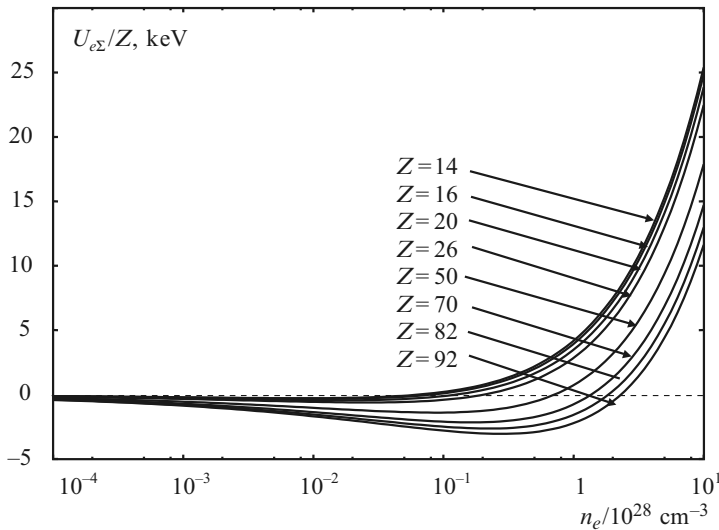


Fig. 11.3. Energy of the degenerate electron gas with low density in the volume of a compressed quasiatom.

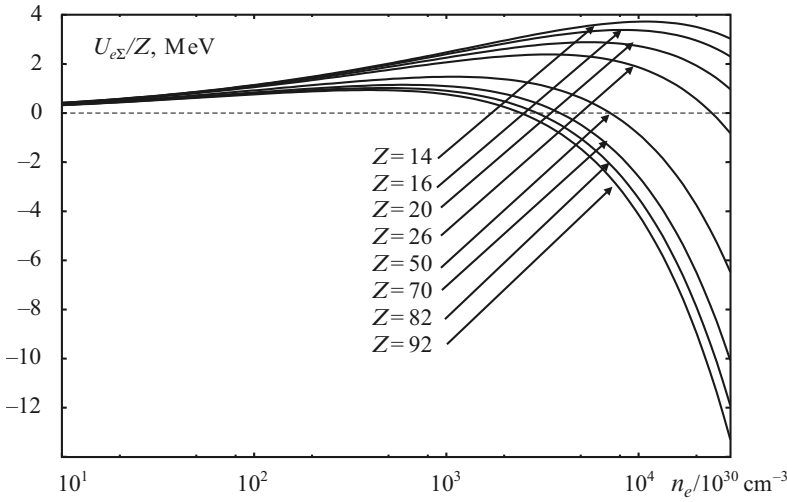


Fig. 11.4. Energy of the degenerate electron gas in the volume of a compressed quasiatome in the region of unstable balance (logarithmic scale).

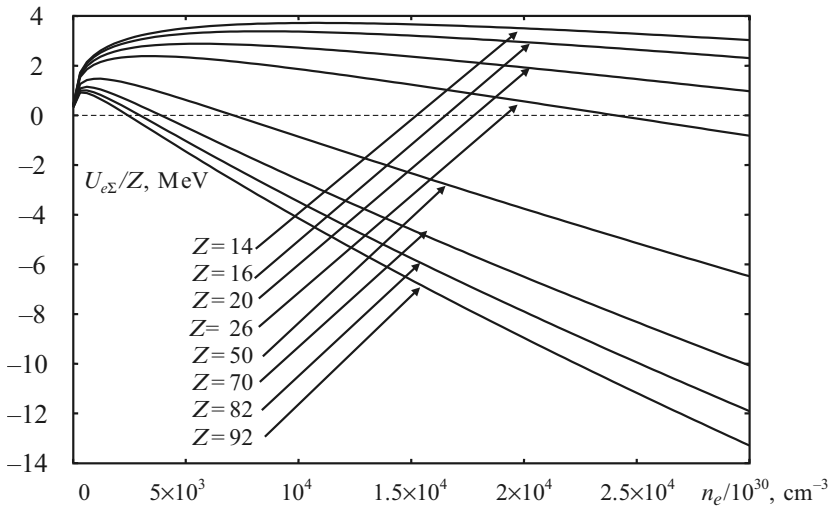


Fig. 11.5. Energy of the degenerate electron gas in the volume of a compressed quasiatome in the region of unstable balance (linear scale).

It follows from Fig. 11.3 that, at a relatively low density of the non-relativistic degenerate gas (i.e., in the interval $n_e \approx (0.3 \text{ to } 3) \times 10^{27} \text{ cm}^{-3}$), each type of nuclei corresponds to a local energy minimum. These minima correspond to the zero pressure, $P_e(Z, n_e) = 0$, on the surface of a compressed quasiatome. The condition $P_e(Z, n_e) = 0$ defines a stable state of the compressed quasiatome which can exist without any external action. It is an

analog (in the region of a high density of electrons) of that metallic type of coupling in crystals which is realized in the degenerate nonrelativistic gas of conduction electrons with a relatively low (on the scale of the problem under consideration) density $n_e \approx 10^{23} \text{ cm}^{-3}$ in metals.

Upon a further compression of the electron-nucleus plasma, there occur the increase in the density n_e and the total and maximum energies of the electron gas and the destruction of the stable states, which corresponds to the transition to the unstable plasma. These dependences define the left half (relative to the maximum) of each plot in Fig. 11.4. To keep such a plasma, it is necessary to apply an additional (negative) external pressure equal to the positive pressure on the compressed quasiautom surface, $P_e(Z, n_e) > 0$.

Upon the further increase in the density n_e [up to $n_{e(cr)} \approx 10^{32}$ to 10^{34} cm^{-3} corresponding to the superrelativistic gas], the total energy of the compressed quasiautom reaches its maximum, which corresponds to the zero pressure $P_e(Z, n_{e(cr)}) = 0$ on the surface of a Wigner–Seitz cell and to the state of unstable balance of the Fermi repulsion forces between electrons and Coulomb attraction of electrons to the nucleus. The structure of the maximum of the mean energy of each electron $U_{e\Sigma}/Z$ is presented in Fig. 11.4 (on the logarithmic scale) and in Fig. 11.5 (on the linear scale).

Upon the further increase in the density of the electron gas, there occurs a decrease in the total energy of the gas of electrons, which corresponds to the negative pressure $P_e(Z, n_e) < 0$ on the surface of a Wigner–Seitz cell. These dependences define the right half (relative to the energy maximum) of each plot in Figs. 11.4 and 11.5. For such a character of the pressure, the system of the compressed quasiautom at $n_{e(c)} > n_{e(cr)}$ is unstable relative to the process of spontaneous self-compression of the plasma and the subsequent unlimited increase in the electron density n_e . In this case, there occurs the compression of the relativistic degenerate gas. This self-strengthening process of irreversible self-compression of the degenerate electron gas is accompanied by the “falling of electron shells of the compressed quasiautom onto a nucleus” and leads to the full collapse of the electron-nucleus plasma in the volume of a specific Wigner–Seitz cell. It is seen from these figures that the mean binding energy of every electron at a sufficiently large electron density can reach many tens of MeV. In this case, the binding energy of the electron-nucleus subsystem of a quasiautom reaches several GeV!

The threshold value $n_{e(cr)} \equiv n_{e(cr)}(Z)$, at which the inversion of pressure and the beginning of the formation of the collapse state occur, depends on the nucleus charge and monotonically increases with decrease in Z , beginning from $n_{e(cr)} \approx 2 \times 10^{32} \text{ cm}^{-3}$ for the heaviest stable nucleus (a nucleus of uranium with $Z = 92$). The dependence of $n_{e(cr)}$ on Z is given in Table 11.1.

Table 11.1. Threshold density $n_{e(cr)}$ and mean threshold energy of electrons necessary for the beginning of the process of Coulomb collapse.

Z	$n_{e(cr)}, \text{ cm}^{-3}$	$U_{e\Sigma}/Z, \text{ MeV}$
92	2.0×10^{32}	0.65
82	3.0×10^{32}	0.75
70	6.0×10^{32}	1.10
50	1.1×10^{33}	1.50
26	3.5×10^{33}	2.40
20	5.8×10^{33}	2.90
16	8.2×10^{33}	3.20
14	1.1×10^{34}	3.70

It seems at the first glance that the necessary values $n_{e(cr)} \geq 10^{32} \text{ cm}^{-3}$ of the critical density of the degenerate gas of electrons are so great that they can hardly be reached under conditions of a terrestrial laboratory. Really, this is not the case.

The point is in that the defining value is not $n_{e(cr)}$, but that additional energy which should be spent on the preliminary compression of the electron-nucleus plasma up to this density.

The performed calculations yield that a relatively small energy is required for the attainment of the threshold of a collapse of a compressed quasiautom [the attainment of the threshold density of electrons $n_{e(cr)}$]. For a nucleus with $Z = 92$, one needs the mean energy $U_{e\Sigma}/Z \approx 0.65 \text{ MeV/electron}$ to reach the threshold density $n_{e(cr)} \approx 2 \times 10^{32} \text{ cm}^{-3}$. This value is much less than that kinetic (Fermi) energy,

$$U_{eF}/Z = (3/4)(3\pi^2)^{1/3} \hbar c n_{e(cr)}^{1/3} \approx 3 \text{ MeV/electron}, \quad (11.42)$$

which should be transferred to the gas of degenerate electrons to reach $n_{e(cr)}$, if we do not account the additional action of the Coulomb attraction of nuclei. The difference of these values, $U_{eF}/Z - U_{e\Sigma}/Z \approx 2.4 \text{ MeV/electron}$, corresponds to a spontaneous increase in the binding energy of electrons with the nucleus in the volume of a compressed quasiautom. The physical explanation of such a favorable (from the viewpoint of a practical realization) effect is related to the following. Upon the additional compression of the degenerate electron gas, the kinetic energy grows (this increases the total energy), and the potential and total energy decrease simultaneously at the

expense of the localization of electrons in the region of the stronger Coulomb field near the nucleus under such a compression.

Thus, about 80 % of all the work which should be performed to reach the threshold of a collapse are carried out by the electron-nucleus system itself at the expense of the Coulomb interaction. For the sake of comparison, we indicate that, without regard for the mentioned work, one can reach the density of the degenerate electron gas

$$n_{e(\text{eff})} \approx \left[(64/81\pi^2)/(\hbar c)^3 \right] U_{e\Sigma}/Z \approx 10^{30} \text{ cm}^{-3} \quad (11.43)$$

if the external forces spent the energy $U_{e\Sigma}/Z \approx 0.65 \text{ MeV/electron}$. Value (Eq. 11.43) is 200 times less than that derived with regard to the electron-nucleus Coulomb interaction.

Respectively, the mean energy $U_{e\Sigma}/Z \approx 0.75 \text{ MeV/electron}$ is required for Pb nuclei ($Z = 82$) to reach the critical density $n_{e(cr)} \approx 3 \times 10^{32} \text{ cm}^{-3}$. Without regard for the Coulomb interaction, one can reach the electron gas density

$$n_{e(\text{eff})} \approx \left[(64/81\pi^2)/(\hbar c)^3 \right] U_{e\Sigma}/Z \approx 1.2 \times 10^{30} \text{ cm}^{-3} \quad (11.44)$$

with such a mean energy.

The dependence of the critical density on the nucleus charge leads to a certain sequence for the transition of nuclei to the state of the formation of a collapse. Upon a gradual increase in the external pressure, electrons from the volumes of the Wigner–Seitz cells with nuclei with maximum charge at their centers will be the first to pass into the state of collapse. If the degenerate electron plasma contains different nuclei in its volume, only those nuclei, for which the condition of the pressure inversion at the corresponding critical density $n_{e(cr)}(Z)$ is satisfied, will pass into the state of electron-nucleus collapse. Thus, the volume of the relativistic degenerate electron-nucleus plasma can simultaneously contain heavy nuclei in the state of electron-nucleus collapse and lighter nuclei in the state with compressed, but not collapsed quasiatoms.

11.1.4. Synthesis of Superheavy Nuclei and Formation of a Nuclear Cluster

Let us consider the influence of the Coulomb collapse of an electron-nucleus systems on the evolution of a nuclear subsystem.

The process of formation of the electron-nucleus collapse starts after the attainment of the critical electron density $n_{e(cr)}$ and leads to the avalanche-type increase in the density of electrons c in the degenerate relativistic gas up to a value comparable to the density of protons n_p in

a nucleus. In this process, there occurs a very significant increase in the binding energy of electrons participating in the electron-nucleus collapse, which changes the character of intranuclear processes and the further evolution of the collapse.

First of all, we note that the evolution of the collapse itself should be studied basing on the total energy of a compressed quasiautom that includes the nucleus energy E_n (Eq. 11.4) and the total energy of electrons $U_{e\Sigma}$ (Eq. 11.35). In this case, the formula for the nucleus energy must be changed with regard to the partial screening of the charge of protons in the nucleus by the degenerate gas of relativistic electrons, which is equivalent to a decrease in the charge of every proton ($e \rightarrow e[1 - n_e/n_p]$) or in the number of protons ($Z \rightarrow [1 - n_e/n_p]Z$).

Figure 11.6 shows the dependence of the total energy of a compressed quasiautom normed per nucleon (a Wigner–Seitz cell in the phase of the formation of the electron-nucleus collapse)

$$\begin{aligned} \Delta E_{en}/A = & \varepsilon_2/A^{1/3} + \varepsilon_3(1 - n_e/n_p)^2 Z_{\text{opt}}^2/A^{4/3} \\ & + \varepsilon_4[1/2 - (1 - n_e/n_p)Z_{\text{opt}}/A]^2 + \varepsilon_5/A^{7/4} + U_{e\Sigma}(Z_{\text{opt}})/A \end{aligned} \quad (11.45)$$

on the atomic number A and the density of the degenerate relativistic electron gas.

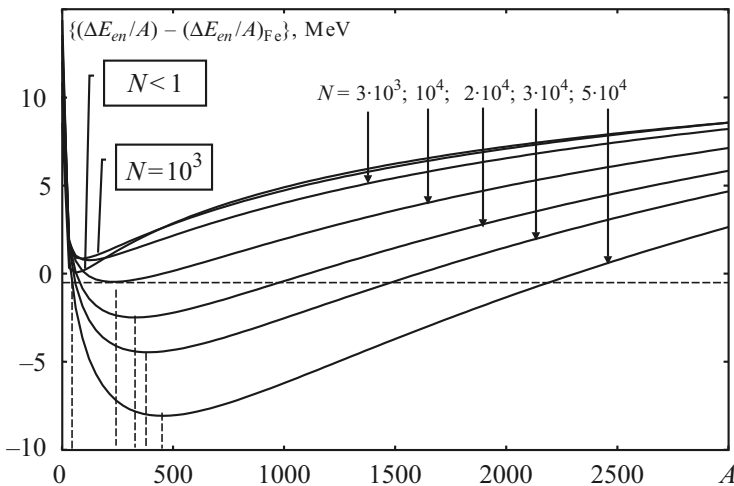


Fig. 11.6. Energy of a compressed quasiautom in the state of partial collapse of the degenerate electron-nucleus plasma with electron density $n_e = 5 \times 10^{34} \text{ cm}^{-3}$. The calculation corresponds to the “classical” line of stability $Z = Z_{\text{opt}}(A)$.

The calculations correspond to the classical line of stability $Z_{\text{opt}}(A)$ for the nucleus charge with regard to screening:

$$\begin{aligned} Z_{\text{opt}} &= A / \left\{ 2 \left[1 + (\varepsilon_3/\varepsilon_4)(1 - n_e/n_p)^2 A^{2/3} \right] \right\} \\ &\approx A / \left[2 + 0.0155(1 - n_e/n_p)^2 A^{2/3} \right]. \end{aligned} \quad (11.46)$$

For convenience of the analysis, the quantity $\Delta E_{en}/A$ was normed by amplitude on the energy of the most stable nuclei of the type of iron and on the invariable total rest energy of all nucleons and electrons, as well as on the binding energy of all nucleons, $\varepsilon_1 A$.

The results of calculations yield that the increase in the density of electrons in the volume of a nucleus is accompanied by a continuous decrease in the energy of the electron-nucleus system (caused by a corresponding increase in its binding energy) and by a displacement of the minimum of this energy to the region of greater masses A .

At $n_e < 10^{30} \text{ cm}^{-3}$ (at $N_e \equiv n_e 10^{30} \text{ cm}^{-3} < 1$), the minimum of the energy of the electron-nucleus system corresponds to nuclei with mass number $A_{\text{opt}} \approx 60$. With increase in the density of electrons n_e , the position of the minimum A_{opt} monotonically increases (see Table 11.2).

Table 11.2. Influence of the electron density on a displacement of the optimum mass number of a nucleus, A_{opt} , characterizing the position of a minimum of the energy of an electron-nucleus system in the compressed electron gas.

$n_e, \text{ cm}^{-3}$	A_{opt}
$< 10^{30}$	≈ 60
10^{33}	≈ 90
3×10^{33}	≈ 130
10^{34}	≈ 250
2×10^{34}	≈ 330
3×10^{34}	≈ 390
5×10^{34}	≈ 450

A displacement of the position of A_{opt} for the minimum of the total energy $\Delta E_{en}/A$ upon a stage-by-stage increase in the electron gas density is accompanied first by the increase in $\Delta E_{en}/A$ from the initial value $\Delta E_{en}(A_{\text{opt}})/A_{\text{opt}} = 0$ at $n_e \ll 10^{30} \text{ cm}^{-3}$ (which corresponds to nuclei of the group of iron) to $\Delta E_{en}(A_{\text{opt}})/A_{\text{opt}} \approx 0.85 \text{ MeV/nucleon}$ at $n_e = 3 \times 10^{33} \text{ cm}^{-3}$.

After the attainment of a local maximum of $\Delta E_{en}(A_{\text{opt}})/A_{\text{opt}}$, there occurs its irreversible decrease: $\Delta E_{en}(A_{\text{opt}})/A_{\text{opt}} \approx 0$ at $n_e \approx 5 \times 10^{33} \text{ cm}^{-3}$ and $\Delta E_{en}(A_{\text{opt}})/A_{\text{opt}} \approx -8 \text{ MeV/nucleon}$ at $n_e = 5 \times 10^{34} \text{ cm}^{-3}$.

The derived values of A_{opt} and $\Delta E_{en}(A_{\text{opt}})/A_{\text{opt}}$ characterize the direct influence of degenerate electrons (at the expense of the own energy of the electron condensate).

Such displacements of the position and depth of the energy minimum for the electron-nucleus system open a way to the formation of extremely heavy nuclei at the expense of the multiple absorption of “ordinary” nuclei by the nucleus with $A \ll A_{\text{opt}}$, being on the collapse formation stage.

We note that it is possible to reach a greater displacement of both the minimum of the total energy A_{opt} and $E_{en}(A_{\text{opt}})/A_{\text{opt}}$. This is conditioned by that the classical line of stability $Z_{\text{opt}}(A)$ for the nucleus charge (Eq. 11.46) corresponds to the condition of the optimization of nuclear parameters on the basis of the formula $d(E_n/A)/dZ = 0$. If we carry out the optimization with regard to the nuclear and electron parameters, we can find the corrected value $Z_{\text{opt}}(A, n_e)$ from the condition $d(\Delta E_{en}/A)/dZ = 0$ which corresponds to a deeper minimum of the electron-nucleus system energy.

In Fig. 11.7, we present the dependence of the energy $\Delta E_{en}/A$ (Eq. 11.45) calculated for one of the versions of the modified relation between the nucleus charge Z and the mass number A ,

$$\begin{aligned} Z_{\text{opt}} &= Ak / \left\{ 2 \left[1 + (\varepsilon_3/\varepsilon_4)(1 - n_e/n_p)^2 A^{2/3} \right] \right\} \\ &\approx Ak / \left[2 + 0.0155(1 - n_e/n_p)^2 A^{2/3} \right], \end{aligned} \quad (11.47)$$

for different values $0.5 \leq k \leq 2$ (here, $Z < A$ in all cases). This condition is far, naturally, from the optimum one.

Equation 11.47 is analyzed for one specific value of the electron gas density $n_e = 3 \times 10^{34} \text{ cm}^{-3}$. The relevant data with the use of the “standard” condition of the line of stability (Eq. 11.46) are given in Fig. 11.6.

Figure 11.7 shows the significant displacements of both the position of the normed energy minimum for the electron-nucleus system in a Wigner–Seitz cell $\Delta E_{en}/A$ and the depth of this minimum even for such simple type of optimization. The position and depth of the energy minimum are shifted from the values $A_{\text{opt}} = 450$, $Z_{\text{opt}} = 156$, and $\Delta E_{en}/A_{\text{opt}} \approx -8 \text{ MeV/nucleon}$ for the standard line of stability (Eq. 11.46) to $A_{\text{opt}} = 1450$, $Z_{\text{opt}} = 650$, and $\Delta E_{en}/A_{\text{opt}} \approx -27 \text{ MeV/nucleon}$ upon the use of relation (Eq. 11.47) between the nucleus charge and the mass number $Z(A, k)$ at $k = 2$.

It is necessary to note that the Coulomb collapse dynamics developing with increase in the masses of nuclei was considered on the basis of the initial premise about that the screening effects can be neglected. Moreover, we took

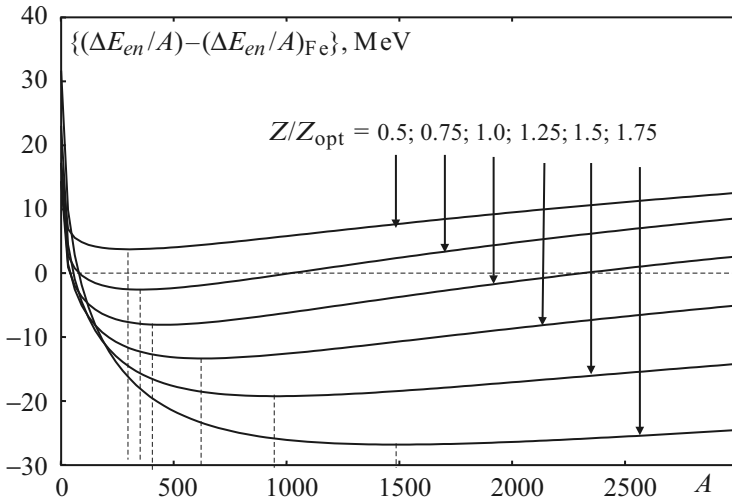


Fig. 11.7. Energy of a compressed quasiautom in the state of partial collapse of the degenerate electron-nucleus plasma with the electron density $n_e = 5 \times 10^{34} \text{ cm}^{-3}$. The calculation corresponds to different values of the deviation of $Z = kZ_{\text{opt}}(A)$ from the “classical” line of stability $Z = Z_{\text{opt}}(A)$; $k = 0.5, 0.75, 1.0, 1.25, 1.5, 1.75$.

a homogeneous electron density $n_e = Z/(4\pi R^3/3)$ in calculations. These assumptions are justified if the screening length R_s exceeds the radius of a compressed quasiautom $R = Z^{1/3}(3/4\pi n_e)^{1/3}$. According to the estimation of A. Migdal, the effect of electron screening in a compressed degenerate relativistic gas of electrons is insignificant if $Z \ll (\hbar c/e^2)^{3/2} \approx 1600$. Thus, we may expect that $Z_{\text{max}} \approx 500$ to 800 is that maximum charge which can be considered on the basis of the above-presented calculations.

The same limitation follows from the results given in Refs. 46, 47. According to this conclusion, the influence of correlation effects in the extremely dense degenerate electron gas described with the static dielectric function $\epsilon_e(k, 0)$ is revealed only in the case where the screening length R_s is considerably less than the quasiautom radius $R = Z^{1/3}(3/4\pi n_e)^{1/3}$ (i.e., at $Z \gg 200$).

In the analysis of the process of collapse for the nuclei with charge $Z > 1000$, it is necessary to consider the effects of screening of the Coulomb field of a nucleus, which will be performed elsewhere.

The synthesis of superheavy nuclei is also favored by that the electric field of a nucleus on the boundary of a quasiautom being compressed (on the boundary of a Wigner–Seitz cell) is completely neutralized by condensed electrons, which strongly weakens the influence of the Coulomb barrier on the internuclear interaction.

Below, we consider this question in more details.

As well known, the cross section of a nuclear reaction,

$$\sigma = (S(E)/E)P(E), \quad (11.48)$$

is defined by the product of the “internal nuclear cross section” σ_0 and the quasiclassical probability of penetration through a Coulomb barrier $V(r)$ (“the Gamow factor”)

$$P = \exp[-2W(E)], \quad (11.49)$$

$$W(E) = \int_{r_1}^{r_2} \sqrt{2\mu(V(r) - E)} dr / \hbar = \sqrt{2\mu\langle V(E) \rangle} |r_2 - r_1| / \hbar. \quad (11.50)$$

In these expressions, r_1 and r_2 are the classical “turning points” upon the motion of one of the particles in the field of the other; μ is the reduced mass of interacting particles, $\langle V(E) \rangle$ is the mean height of the potential barrier lying above the level E . Upon the interaction of a particle with the screened nucleus, the minimum and maximum values of r_1 correspond, respectively, to the nucleus radius R and the screening radius R_s .

The formula for the reaction cross section yields two obvious ways to increase the barrier transparency and to attain a high efficiency of the synthesis: the increase in the energy or temperature of particles (this is the region of thermonuclear or accelerator-based synthesis) and the decrease in the barrier width at low temperatures (this is the pycnonuclear synthesis upon the compression of the system, as well as the muonic catalysis upon the decrease in the screening radius at the expense of the capture of a heavy μ -meson by one of the nuclei participating in the synthesis).

Muonic catalysis at low temperatures was proposed independently by Frank (see Ref. 51) and Sakharov (Ref. 59) and then was verified in experiments (see Ref. 50). The complexity to use such attractive model is related, in the first turn, to that the heavy μ -mesons screening the field of a nucleus are unstable particles and have lifetime of about c , which is insufficient for practical applications.

Pycnonuclear synthesis was proposed in works of Zel'dovich and Gershtein (1960) and Harrison (1964). To satisfy the conditions for the energy-gained synthesis on the basis of the nuclei of hydrogen under conditions of a terrestrial laboratory, it is necessary to solve the problem of compression of the condensed medium by 10^3 to 10^4 times, which can lead to the decrease in the barrier width by 10 to 15 times. In the case of the use of heavier nuclei, the required degree of compression of targets sharply increases. The experience of the recent 50 years has shown that the use of external drivers for these purposes (pulse magnetic fields, laser beams,

and ion beams) does not allow one to reach the required compression: the maximum compression attained in the best experiments did not exceed 100, which is clearly insufficient for practical applications.

The very intense search for the optimum synthesis led to that all the leading laboratories over the world chose the thermonuclear mechanism of synthesis as the main one at the beginning of the 1950s. At those laboratories, the researchers are studying the process of synthesis of the nuclei of deuterium and tritium in quasistationary systems for the long-term keeping of a high-temperature low-pressure plasma in magnetic traps or by using a short-term compression (by at least 100 to 1000 times) and an inertial keeping under the uniform laser- or beam-involved compression of a condensed target.

The experience of the last 50 years has demonstrated that the way to realize nuclear reactions which is based on the thermonuclear mechanism turns out to be extremely expensive and inefficient under terrestrial conditions (as distinct from stars, where the stabilization of a hot thermonuclear plasma is realized naturally and efficiently by gravitational forces).

The study of the synthesis on the basis of the inertial keeping of a hot plasma with the use of a laser or beam driver led also to the conclusion about the low efficiency of such a method of optimization.

Contrary to those inefficient ways to provide the synthesis of nuclei, the way proposed by us is based on the use of a distinctive internal driver, whose role can be played, along with the other mechanisms, by the process of formation of the Coulomb collapse.

Consider some of the main details of the process of synthesis in a compressible medium.

It is seen from Eq. 11.50 that the barrier width $|r_2 - r_1| = R_s - R$ (which is a linear function of the coordinate) affects the Gamow factor more considerably than its mean height $\langle V(E) \rangle$ appearing in $W(E)$ under the square root symbol. This circumstance leads to the very sharp decrease in the influence of a Coulomb barrier on the synthesis of nuclei upon the decrease in the radius of a compressed quasihomogeneous atom under its collapse (which is equivalent to a decrease in the barrier width).

Upon a moderately strong compression of the degenerate relativistic gas of electrons (up to the density $n_{e(cr)} \approx 10^{32} \text{ cm}^{-3} < n_e \ll n_p \approx 10^{38} \text{ cm}^{-3}$ which is significantly less than the intranuclear density of protons n_p), the width of the Coulomb barrier sharply decreases: from the initial value equal to the screening radius $R_{s(0)} \approx 0.885(\hbar^2/m_e e^2)/Z^{1/3}$ in atoms and ions with a low degree of ionization or in a noncompressed degenerate electron plasma up to the value $R_s \approx R = (3Z/4\pi n_e)^{1/3}$ defined by the radius R of a Wigner–Seitz cell. In particular, for a nucleus with $Z = 92$,

the initial screening radius $R_{s(0)} \approx 0.1 \rho A$. The width of the Coulomb barrier at the density $n_{e(cr)} \approx 10^{32} \text{ cm}^{-3}$ decreases by $R_{s(0)}/R_s \approx 16.6$ times, at the density $n_e \approx 10^{33} \text{ cm}^{-3}$ by $R_{s(0)}/R_s \approx 36$ times, at the density $n_e \approx 10^{35} \text{ cm}^{-3}$ by $R_{s(0)}/R_s \approx 166$ times, and, finally, at $n_e \approx 10^{37} \text{ cm}^{-3}$ by $R_{s(0)}/R_s \approx 775$ times. Since the probability of the tunnel passage of nuclei through a Coulomb barrier in the compressed electron-nucleus plasma depends mainly on the barrier width, such a cardinal decrease in the barrier width increases sharply the probability of the synthesis.

Below, we present other estimations demonstrating the efficiency of the synthesis.

The nuclei of a cold compressed electron-nucleus plasma can be characterized by the effective energy of relative motion T_{eff} which should be possessed by the same, but free nuclei in order to approach one another at the same distance R_s . For the purely Coulomb nonscreened interaction of two nuclei with charges Z_1 and Z_2 (this corresponds to the considered case of a relativistic degenerate electron gas), this energy is defined as $T_{\text{eff}} \approx Z_1 Z_2 e^2 / R_s$.

For example, a free nucleus of Pb and a proton approach each other up to the distance $R_s \approx 1.2 \times 10^{-12} \text{ cm}$ (this corresponds to the density $n_e \approx 10^{36} \text{ cm}^{-3}$) if they have the relative energy $T_{\text{eff}} \approx 15 \text{ MeV}$. Respectively, the interaction of two nuclei of Pb in the composition of the degenerate electron gas with the same density corresponds to the energy $T_{\text{eff}} \approx 1400 \text{ MeV}$.

With increase in the density of electrons, T_{eff} is also increased, and the probability of the tunnel effect becomes greater.

We indicate one more circumstance. The process of synthesis with the absorption of nuclei of the target leads to the appearance of a “superfluous” binding energy which will be released in various channels (γ -emission, neutron emission, emission of nuclear fragments, etc.) and will cool the nucleus. One of the channels is related to the creation of various “normal” nuclei and to their emission from the volume of the growing superheavy nucleus. For example, after the absorption of several nuclei of the target with $A_T \approx 1$ to 200 for a short time, the great excess of binding energy can lead to the emission of several light nuclei with $A_L \approx A_T$ by one heavy nucleus with $A_H \approx 300$ to $500 > A_T$.

The emission of synthesized nuclei and nuclear fragments is a competitive channel relative to the other channels of cooling the nucleus. In this case, the greater probabilities of the creation and emission are possessed by ordinary even-even nuclei (which are similar to α -particles and isotopes like $\text{C}^{12}, \text{O}^{16}, \dots, \text{Pb}^{208}$) already existent in the ground (nonradioactive) state in the volume of the superheavy nucleus. In fact, every superheavy nucleus is a “distinctive microreactor” for the transmutation of the “ordinary” nuclei

of the target into various combinations of nucleons. In this case, the superheavy nucleus mass also gradually increases. The process of transmutation is completed after the utilization of all nuclei or after the evolution of the superheavy nucleus to a stable state with $A_{\max}$.

The astrophysical aspects of these processes will be considered in Sec. 11.1.5.

As a result of the same tunnel effect, a greater probability is characteristic of the processes of tunnel emission of nuclear fragments with small charge upon the decay of nuclei in the process of synthesis at the expense of the absorption of heavy nuclei (if such a decay is energy-gained).

It is obvious that the greatest probability is characteristic of the emission of such maximally stable nuclei which are formed through fluctuations (in the form of clusters) in the volume of the growing superheavy nucleus and have a greater probability "to approach" the surface of the nucleus and not to decay. Therefore, the binding energy of a cluster must be significantly greater than that corresponding to the temperature of nucleons in the superheavy nucleus.

At the same time, the clusters with a small binding energy (just they correspond to unstable, i.e., radioactive, nuclei after the escape from the maternal nucleus) decay in the volume of a big superheavy nucleus at once after their creation. Such clusters appearing through fluctuations and decaying at once ("flickering" clusters) are characterized by a very small probability to approach the surface nucleus in the bound form and to leave it. For this reason, no radioactive nuclei are observed in practice in the process of transmutation of superheavy nuclei.

This situation basically differs from the spontaneous decay of "ordinary" (let even heavy, but not transuranium) nuclei.

Due to a relatively small size of "ordinary" nuclei, the probability for an unstable nucleus formed through fluctuations to leave the maternal nucleus without decay inside it increases abruptly. For superheavy nuclei, such a process has a low probability.

One of the main reasons for such a difference is obvious and is related to the following. With increase in the atomic number A of a superheavy nucleus, the relative share of the near-surface nucleons with respect to all nucleons of the nucleus continuously decreases (proportionally to the ratio of the nucleus surface and its volume, i.e., inversely proportionally to $A^{1/3}$). The weight share of those clusters with low binding energy which can be formed in the near-surface layer and then can leave the superheavy nucleus as radioactive products decreases to the same extent. Contrary to this case, clusters with high binding energy can reach the surface and leave the nucleus even in the case where they were formed in the inner part of the superheavy nucleus.

The process of self-cooling of a superheavy nuclei under consideration is very similar to the emission of stable α -particles from the volume of radioactive nuclei of the uranium–thorium series. This process is also a direct analog of the self-cooling of a liquid upon the evaporation of stable molecules. During such a cooling, unstable molecular groups including several weakly bound molecules leave a liquid very seldom.

One more reason for the phenomenon of the “nonescape” of radioactive nuclear clusters (in fact, of daughter nuclei) from a growing superheavy nucleus can be related to the fact that, under conditions of the very strong Coulomb interaction of the protons of the nucleus and the compressed degenerate electron gas surrounding this nucleus, the nucleus surface becomes more diffuse (the sharp boundary of the nucleus disappears). For this reason, the escape of a formed nuclear cluster from a maternal nucleus and its associated interaction with surface nucleons of this nucleus become close to an adiabatic process. Whereas, for the “ordinary” nuclei with a clearly pronounced boundary, such an escape corresponds almost always to a pulse (jump-like) process.

Such a difference in the processes of interaction leads to different final states of a daughter nucleus. In the case of a sharp boundary of the maternal nucleus, the probability of the escape of a nonequilibrium excited daughter nucleus is large, but it abruptly decreases in the presence of a diffuse boundary.

These results directly follow from the general principles of the evolution of any resonance system in the presence of an external perturbation. It is obvious that, under an adiabatic perturbation, the probability of the formation of a final nonequilibrium state of such a system is extremely low. The formal aspect of this result is related to that only low frequencies are present in the spectrum of an adiabatic perturbation, and the process of excitation of nuclei corresponds to very high frequencies.

Starting from such arguments, it becomes obvious that stable nuclei will leave the superheavy nucleus with highest probability.

The presented scenario of the collective synthesis is basically different from the commonly accepted one of the natural nucleosynthesis, where isotopes and elements are formed in the individual random s - and r -processes with absorption of neutrons. In such random and nonoptimized (relative to the requirement for the binding energy to be maximum) processes, radioactive nuclei will be inevitably created.

In conclusion, we note that the question about the limiting degree of self-compression of the system of degenerate relativistic electrons in a compressed quasiautom (i.e., about both the achievement of the maximum ratio n_e/n_p and the possibility to weaken the interaction of protons in a

nucleus) can be solved only within the full analysis of all the variety of electron-nucleus transformations (including the processes of protonization and neutronization and the creation of electron-positron pairs) for specific nuclei. A part of these questions will be considered below in Sec. 11.1.6. At the same time, we can surely state that the limiting density of the compressed electron gas, $n_{e(\max)}$, will be always less than the density of protons in a nucleus, n_p , because the single reason for the compression of a relativistic degenerate gas of electrons is its interaction with the efficient nucleus charge $Z^* = Z(n_p - n_e)/n_p$ decreasing upon compression.

The considered self-similar mechanism of the formation of the state of electron-nucleus collapse has a threshold and requires a preliminary attainment of the specific energy of a compressed degenerate electron plasma at a level of about 0.65 MeV/electron, which corresponds to a high critical density of the degenerate gas of electrons $n_{e(cr)}$. This density is very high as compared to the typical value $n_e \approx 10^{23}$ to 10^{24} cm $^{-3}$ characteristic of, for example, the natural state of the conduction zone in metals. At the same time, the value $n_{e(cr)}$ is by many orders less than that final density $n_e \approx 10^{36}$ to 10^{37} cm $^{-3}$ which is automatically reached after a preliminary compression in the self-controlled process of formation of the state of collapse and is comparable to the intranuclear density of nucleons n_n . In the state of collapse, the directivity and the probability of nuclear reactions are basically changed, the influence of the Coulomb barrier sharply decreases, and the spontaneous synthesis of superheavy stable nuclei along with the accompanied synthesis of lighter nuclei become possible. The considered method of combined nuclear transformations (fusion-fission processes) on the basis of the self-controlled electron-nucleus collapse is basically different from the “force” method of synthesis of nuclei which requires a continuous force action.

Very important is the circumstance that the process of formation of the state of electron-nucleus collapse occurs with the participation of the own electron system. For this reason, this process is maximally adapted with respect to the nucleus and runs as a completely consistent process.

The consequences of the formation of such a collapse are numerous and obvious. A part of them is related to the possibility of the synthesis of superheavy nuclei with masses exceeding the mass of transuranium isotopes by many times. Such a synthesis becomes possible because the energy minimum of the electron-nucleus collapse in the process of its rapid formation moves from the region of mass numbers $A_{\text{opt}} \approx 60$ (characteristic of ordinary nuclei and a low density of electrons) to $A_{\text{opt}} > 2000$ to 3000 (at $n_e \approx 10^{35}$ to 10^{36} cm $^{-3}$). This displacement ensures the energy gain of the process of synthesis upon the mutual absorption of nuclei and establishes the

basic premise for the formation of stable superheavy nuclei. The high efficiency of the very synthesis is defined by the circumstance that the Coulomb repulsion for nuclei in the state of electron-nucleus collapse turns out to be strongly weakened due to the strong screening of the nucleus field by the degenerate and compressed relativistic electron gas. This circumstance can lead to a specific selection of those types of nuclei, for which the interaction will be characterized by a greater probability. On the one hand, the beginning of the formation of a collapse is significantly facilitated for heavy initial nuclei with large charge Z . On the other hand, the interaction and subsequent synthesis with the participation of only the nuclei with a large charge Z on the stage far from the limiting compression (when the size of a Wigner–Seitz cell approaches the size of a nucleus) can be essentially hampered by the residual Coulomb repulsion. In such a situation, optimum can be a target, in which the small number of initial heavy nuclei is rarefied by the large number of extremely light nuclei. In this case, each type of nuclei fulfils a certain function: on the basis of heavy nuclei, extremely compressed electron-nucleus systems of the type of a Wigner–Seitz cell are formed in the state of collapse, whereas light nuclei, whose charge is insufficient for the collapse formation, are efficiently absorbed by compressed electron-nucleus systems and increase the masses of the latter, by playing the role of a “constructional material”.

It is necessary to note that such an idealized model describes only one of the possible variants of the optimum composition of a target. There exist many factors which can lead to a completely different optimum composition. Such factors include, in particular, the degree of the limiting compression of a Wigner–Seitz cell and the anomalous radial distribution of near-surface nucleons in a heavy nucleus, which can occur due to a very strong interaction of these nucleons with the extremely compressed degenerate electron gas. These questions need the further detailed studies.

In the scope of the region of a formed electron-nucleus collapse, the interaction of “normal” nuclei with a compressed quasiatome can run with the release of a very large binding energy. For this reason, they can be characterized by a great variety including a combination of numerous acts of synthesis and fission with the formation of new nuclei and the accompanied proton–neutron transmutations, though the general tendency corresponds to the evolution toward the attainment of a minimum energy at A_{opt} .

In particular, the great binding energy, which is released at every act of absorption of one or several heavy nuclei by a compressed quasiatome in the process of formation of the electron-nucleus collapse, can induce the simultaneous emission of several lighter nuclei or one superheavy nucleus into the environment. Such a mechanism of “recycling” of nuclei of a target

into other nuclei was first proposed in Ref. 52. Apparently, the synthesis of a wide spectrum of superheavy nuclei and “normal” nuclei with anomalous isotope ratios, which was observed after the action of a pulse coherent driver on chemically pure targets in the experiments performed at the Electrodynamics Laboratory “Proton-21”, is related just to such processes.

The other obvious consequence of the process of formation of the state of electron-nucleus collapse is the fact that it can ensure, in fact, the transition from the region of “normal” nuclei to the region of such superheavy nuclei, where the stability of nuclear matter is ensured by other mechanisms (in particular, by the pion condensate).

It is probable that the effect of a displacement of the minimum of the total energy of the electron-nucleus system ΔE_{en} upon the attainment of the critical electron density $n_{e(cr)}(Z)$ can be namely that key mechanism which allows both the realization of the direct synthesis of superheavy nuclei with charges in the scope of region 2 with $100 < Z < 1600$ (Fig. 11.1) and the rapid transition from region 1 of stable “normal” nuclei with $A < 240$ and $Z \leq 92$ to region 3 of the nuclei with a self-formed pion condensate and with $Z > 1600$ (Fig. 11.8). For the sake of justice, we note that the above-considered mechanism of the interaction of nuclei with the degenerate electron gas characterizes quite adequately the nuclei

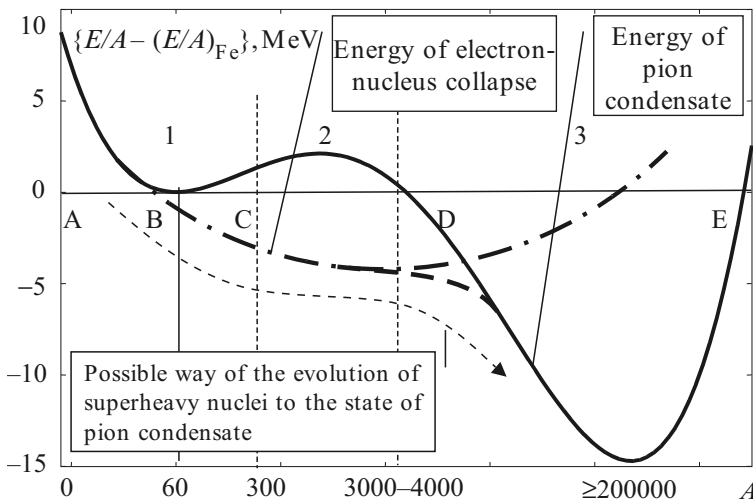


Fig. 11.8. Assumed scheme of the dependence of the nucleus energy at the normal electron and nuclear densities (upper curve in region ABC), in the presence of an electron-nucleus collapse (lower curve in region BCD), and in the presence of a pion condensate in the nucleus volume (lower curve in region DE).

with $Z \leq 600$, because correlation effects for such nuclei in the volume of the superrelativistic electron gas are insignificant. In the region $600 < Z < 1600$, such effects should be considered, which will be made elsewhere.

In conclusion, it is worth noting that the problems related to the character and the directivity of nuclear reactions in the electron-nucleus collapse zone are more complex as compared to analogous problems for ordinary nuclei. Moreover, the solution of the former involves the processes of the type of protonization and neutronization of nuclei with the simultaneous participation of electrons and positrons, as well as the processes of absorption and creation of nuclear fragments. A considerable part of these effects will be considered below in Sec. 11.1.6 To this circle of problems, we refer the question about the limiting degree of self-compression of the system of degenerate relativistic electrons in a compressed quasi-atom. The question about the stability of such superheavy nuclei formed in the process of collapse of the electron-nucleus system, which includes the nucleus and the degenerate electron Fermi-condensate, should be solved with regard to the whole totality of nuclear transformations (including nuclear r -processes, the nuclear decay processes competitive with the synthesis, and the pion condensation in the volume of synthesized superheavy nuclei). These problems are extraordinarily complex and will be considered in the subsequent works.

11.1.5. Mechanism of the Nucleosynthesis of Superheavy and Neutron-Deficient Nuclei upon the Sequential Action of the Gravitational and Coulomb Collapses in Astrophysical Objects

The Question about and Unsolved Problems of the Natural Nucleosynthesis in the Universe. The above-considered conditions of the formation of the electron-nucleus collapse require to realize the threshold degree of a compression of matter (up to the state of degenerate relativistic electron gas). In this case, the threshold degree of a compression will be minimum for nuclei with the largest charge.

Such conditions can be satisfied, in particular, in astrophysical objects such as the crust and the core of a neutron star and, under certain additional conditions, at an ordinary terrestrial laboratory. The specific conditions for the realization of a collapse under conditions of a terrestrial laboratory without the use of a preliminary supercompression of the target and without regard for the influence of the gravitational forces will be considered in the next chapter.

Now we will analyze the possibility to realize the process of formation of the electron-nucleus collapse in astrophysical objects. The importance of

such a consideration (besides the methodical questions related to the general problems of the Coulomb collapse of an electron-nucleus system) consists in that the use of the general conception of Coulomb collapse allows us to give a conjectural answer to many unsolved problems of astrophysics and, in particular, to the problem of the nucleosynthesis in the universe.

The problem of the nucleosynthesis of chemical elements and isotopes is intensively studied for recent 50 years. At present, commonly accepted is the idea of that all the groups of nuclei were formed at different times at the expense of basically different mechanisms: helium and other light elements with $A < 12$ were formed in the reactions of thermonuclear synthesis at the early stage of the universe evolution; the elements beginning from carbon and to the nuclei of the “iron peak” (Mn, Fe, and Ni) are formed at a high temperature in the sequence of the reactions of thermonuclear synthesis which started from nuclei C^{12} and O^{16} and ran in the process of evolution of stars with large masses; and the elements heavier than iron were formed in the central part of massive stars in the reactions of sequential capture of neutrons at the low (s -processes) and high (r -processes) intensities of the neutron flux.

We mention two groups of nuclei which cannot be formed in any above-discussed way and remain beyond the frame of such logical system.

First of all, these are the nuclei of elements Li, Be, and B which are easily destroyed in thermonuclear reactions due to the insufficiently high binding energy.

Finally, we point out 34 medium and heavy “neutron-deficient nuclei” with a relatively low content of neutrons: Se^{74} , Kr^{78} , Kr^{80} , Sr^{84} , Mo^{92} , Mo^{94} , Ru^{96} , Ru^{98} , Pd^{102} , Cd^{108} , In^{113} , Sn^{112} , Sn^{114} , Sn^{115} , Te^{120} , Xe^{124} , Xe^{126} , Ba^{130} , Ba^{132} , Ce^{136} , Ce^{138} , Sm^{144} , Gd^{152} , Dy^{152} , Dy^{158} , Er^{162} , Er^{164} , Yb^{168} , Hf^{174} , W^{180} , Os^{184} , Pt^{190} , and Hg^{196} . Almost all the neutron-deficient nuclei have an even mass number or a magic number of protons (except for In^{113} and Sn^{115}). Neutron-deficient nuclei cannot be created in the equilibrium s - and r -processes with absorption of neutrons on the slow or fast evolutionary stages of development of stars. It is considered that their origin is directly related to the weak interaction with the participation of neutrinos in nuclear reactions accompanying the gravitational collapse of heavy stars. There exist several serious arguments which cast some doubt on this viewpoint.

We consider the assumption that the neutron-deficient nuclei can be formed in the reactions of protonization of nuclei



at the expense of the interaction of neutrinos emitted by the central region of a collapsing star in the process of neutronization of nuclei with “normal”

nuclei of a dense matter of the internal layer of its shell to be sufficiently controversial with regard to the extremely low cross section of the interaction of neutrinos with any nuclei.

In the central part of a star, the possibility for reaction Eq. 11.51 to run is very limited (in fact, it is forbidden because of the very high Fermi energy, $E_F = (3\pi^2)^{1/3} \hbar c n_e^{1/3} \geq 20$ to 30 MeV, of degenerate relativistic electrons in the composition of the electron-nucleus plasma entering the matter of a star. For reaction Eq. 11.51 to run, the newly created relativistic electrons must have energy exceeding E_F , which can be only under the condition that the initial relativistic neutrinos have energy significantly exceeding E_F . It is difficult to substantiate the existence of a great number of high-energy neutrinos on the basis of the standard model of gravitational collapse (the typical neutrino energy upon a gravitational collapse is equal to 10 to 15 MeV). We also note that because the fluxes of neutrinos and antineutrinos from the region of the gravitational collapse are approximately equal each to other, the neutronization reaction inverse to Eq. 11.51,



will lead to the destruction of such neutron-deficient nuclei. This is favored by the circumstance that, due to the excess of electrons in the composition of the degenerate plasma, reaction Eq. 11.52 is more probable than Eq. 11.51, since the former can run with the participation of both fast and slow antineutrinos.

In the scope of the outer shell of a star, where the Fermi energy E_F decreases sharply because of a lesser concentration of degenerate nonrelativistic electrons, the result remains the same: due to the isotropic character of the motion from the center of a collapsing star, the flux of neutrinos will be very weak, which also makes reaction Eq. 11.51 unlikely.

One more argument of the insufficient conclusiveness of the purely “neutrino-based” scenario for the origin of neutron-deficient elements is related to the fact that, according to reaction Eq. 11.51, the abundance of their isotopes X_{Z+1}^A must be, as a rule, proportional to the abundance of the isotopes of maternal nuclei X_Z^A , which is very frequently violated.

It is necessary to note that all the known scenarios for the creation of medium and heavy nuclei (both normal and neutron-deficient ones) are based only on the single logical scheme: the creation of nuclei occurs by the scheme of the energy-gained synthesis by means of the sequential increase in the mass number upon the capture of neutrons.

Such a scheme has significant and, apparently, insurmountable flaw: if, in the chain of sequential transformations of heavy isotopes, the synthesis is improbable or has low efficiency on its some link (because the created

nuclei decay very rapidly), then this “weak link” interrupts or decelerates the synthesis of all heavier nuclei. As precursors of such “weak links” of the stage-by-stage synthesis, we mention, for example, magic or twice magic nuclei. One of such nuclei is Bi^{209} . After the absorption of a neutron, the formed nucleus Bi^{210} is unstable with respect to the α -decay and is transformed in Pb^{206} .

A number of unsolved questions in the problem of nucleosynthesis on the basis of r -processes is related, for example, to the characteristic peaks in the abundance of elements at $N = 50, 82$, and 126 . It was shown many times (e.g., Ref. 58) that, in order to form such peaks at a single temperature of the medium, the densities of neutrons different by 6 orders ($n_n = 8 \times 10^{19}$, 3×10^{23} , and 10^{26} cm^{-3}) are needed! In addition, it is difficult to point out the sources with a very high density of free neutrons required for the synthesis of heavy nuclei in the r -process.

These fundamental problems, despite the certain successes reached in the comprehension of different special questions (e.g., Ref. 57), remain unsolved up to now.

Below, we consider the basically different scheme of nucleosynthesis running in the shell of a gravitating massive star. This scheme is based on the combination of the reactions of synthesis and fission of heavy and superheavy nuclei upon the Coulomb collapse of the electron-nucleus plasma, the conditions for the realization of which are established in the preceding process of gravitational collapse of a heavy star. Such a sequential tandem of the gravitational and Coulomb collapses allows us not only to explain the most part of element and isotope anomalies, but to answer some questions related to the evolution of stars in the process of gravitational collapse (in particular, with its help, we can clarify the process of fast release of a great energy on the inner surface of the shell of such a star leading to the dispersion of its shell).

Qualitative and Quantitative Comparison of the Conditions for a Realization of the Gravitational Collapse of a Heavy Star and the Coulomb Collapse of a Compressed Atom. Starting from the circle of questions considered in the previous chapter, we advance the assumption that the main mechanism and regularities of the sequential tandem of the gravitational and Coulomb collapses of a degenerate electron-nucleus plasma are based on two independent, but mutually consistent phenomena (two stages):

Stage 1. *The evolution of any star with mass of more than $1.45M_\odot$ terminates, as a rule, by the gravitational collapse of the central iron region of a star.*

Here, $M_{\odot} \approx 2 \times 10^{33}$ g is the sun's mass.

In the process of collapse, the density of a neutral electron-nucleus plasma ρ reaches the value of $\approx 10^{11}$ to 10^{14} g/cm³, whereas the concentration of the degenerate relativistic electron gas in this plasma (with regard to the process of neutronization) reaches $n_{e(\max)} \approx 10^{34}$ to 10^{37} cm⁻³. The process of formation of the gravitational collapse begins after the compression of the neutral electron-nucleus plasma by the gravitational force up to such a density, at which the electron component of the plasma will transit from the state of degenerate nonrelativistic electron gas to the state of relativistic gas. The role of a compressing driver in this process is played by the gravitational field.

The process of gravitational collapse terminates on the hydrodynamic stage with the formation of a neutron star or induces the relativistic collapse with the formation of a black hole. In the process of gravitational collapse, a very great energy ($E_G \approx 10^{54}$ erg) is released. To separate the very heavy outer shell containing about half a mass of the star and then to accelerate it to the velocity $v_S \approx 10^9$ cm/s, the energy $E_S \approx 10^{51}$ erg is required. Though the process of gravitational collapse is well studied on the whole, there is a number of circumstances testifying to the existence of essential drawbacks in the comprehension of specific mechanisms of release and transport of this energy from the collapse region to the shell. It was convincingly shown in many works (see, e.g., Ref. 60) that neither a shock wave, nor the flux of neutrinos can ensure the transfer of the great energy E_S from the central region of a gravitating star to the inner surface of its dispersing shell.

Stage 2. *At the final stage of the gravitational collapse, the conditions for the Coulomb collapse of local regions of the electron-nucleus plasma are satisfied.*

The state and evolution of every part of the compressed electron-nucleus plasma depend on a local density, temperature, and the charges of nuclei Ze in the composition of this plasma. They are defined by the balance of the pressure compressing the plasma due to the mutual Coulomb attraction of electrons and a nucleus in the volume of a neutral Wigner-Seitz cell and the kinetic Fermi pressure of the gas of these electrons hampering its compression. We showed earlier that, for each type of nuclei in the composition of this plasma, there exists the proper threshold concentration $n_{e(cr)}$ of the degenerate relativistic gas of electrons, beginning from which the Coulomb pressure compressing the plasma exceeds the kinetic Fermi pressure. The minimum threshold value of $n_{e(cr)}$ corresponds to the maximum charge of nuclei Z and can be reached at the expense of the action of an external driver. To attain the threshold concentration $n_{e(cr)}$, it is necessary to increase the mean energy of every electron of the degenerate plasma up to

0.65 MeV/electron for nuclei U, 0.75 MeV/electron for Pb, 2 MeV/electron for Fe, and 2.6 MeV/electron for Ca. The attainment of the threshold value of $n_{e(cr)}$ induces the process of Coulomb collapse of the electron-nucleus plasma in the volume of every Wigner–Seitz cell.

It is necessary to note that the energy of the strongly compressed electron-nucleus system, in which electrons and nucleons are the partners with quite equal rights, changes tremendously in the process of Coulomb collapse. The process of collapse leads to a displacement of the position of the binding energy maximum for this system from the ordinary one for uncompressed atoms with mass number $A = 60$ to very high values $A > 1000$. This circumstance opens a way to the synthesis of heavy and superheavy nuclei and to the associated synthesis of any nuclei, including neutron-deficient ones.

The preliminary compression of the medium for the attainment of the threshold of a Coulomb collapse can be created on the earth on the basis of the use of an optimum driver (this mechanism will be analyzed partially in the next chapter). However, it is more natural to use the process of gravitational collapse compressing the electron-nucleus plasma in the volume of a star as a driver. The process of gravitational collapse leads to such densities of degenerate electrons $n_{e(max)}$ which considerably exceed the threshold value $n_{e(cr)}$ for the Coulomb collapse for any nuclei, except for the lightest ones. These results will be considered below in details.

Now we make a qualitative comparison of the gravitational and Coulomb collapses.

First of all, we note that the threshold conditions of a Coulomb collapse (the threshold charge of a nucleus Z_{cr} at the given concentration n_e of the degenerate relativistic gas of electrons or, respectively, the threshold concentration of the degenerate relativistic gas of electrons $n_{e(cr)}$ depending on the nucleus charge) remind the analogous conditions for a gravitational collapse (the critical star mass equal to $1.45M_\odot$ and the state of a shell in the form of a degenerate relativistic gas of electrons). There are certain differences: the condition for a gravitational collapse is toughly connected with the fixed critical mass of a star and does not depend on the density of a degenerate relativistic gas of electrons, whereas the Coulomb collapse can run for any nucleus charge but, in this case, requires the threshold density of the gas of electrons consistent with the nucleus charge.

We also note that the initial stages of the process of attraction of the shell of a heavy star to its center and, respectively, the process of attraction of electrons to the nucleus are described by the analogous interaction energies depending on the distance as ($\sim 1/r$). Both these processes become especially similar, when we consider the shell of a star and the electron shell

of a compressed atom in the form of a degenerate Fermi-gas: the shell of a star is a neutral plasma which is degenerate in the electron component and is attracted to the center at the expense of the gravitational forces; whereas a compressed atom in the form of a Wigner–Seitz cell includes a degenerate electron gas attracting to the nucleus. In both systems, the attraction to the center is a stimulus to the compression, whereas the kinetic Fermi pressure of the degenerate gas of electrons hampers the compression. Of course, there are the considerable differences in these processes (the Coulomb field of a nucleus can be screened by electrons, but the gravitational field cannot be screened; a compressed atom has a separated central core, but such a center is absent in a star), but these differences can affect only the collapse threshold and its final stage, but not the very existence of a collapse.

In view of the qualitative similarity of these systems, it would be expected that the ways of their evolution are also similar. However, we observe basically different patterns of their evolution in practice.

As known, if the premises of the gravitational collapse of the shell of a heavy star are satisfied and the Fermi pressure cannot keep the shell of a star from the catastrophic gravitational self-compression, then the gravitational collapse is the almost inevitable final stage of the star's evolution. Therefore, a quite natural question arises: why was an analogous Coulomb collapse of a compressed atom not earlier even discussed?

At the first glance, the reason is seemed to be in that one considers different systems: an atom is a purely quantum object, whereas a star is a classical one. However, it is easy to verify that a star, whose shell consists of a degenerate neutral electron gas, is also a purely quantum object like an atom.

To answer this question, first we consider both these processes in a simplified form.

First, we will analyze the conditions for the self-compression of a star at the expense of the action of the gravitational forces.

As known, the state of a stellar matter prior to the gravitational collapse corresponds to a degenerate electron gas, in the volume of which the nuclei occupy the positions close to an ordered cubic lattice. This lattice is formed from adjacent neutral Wigner–Seitz cells.

The energy of the gravitational interaction of a stellar matter is

$$W_G = -gGM^2/R = -gG(kM_n)^2(4\pi/3)^{1/3}(NZ)^{5/3}n_e^{1/3}, \quad (11.53)$$

where $g \approx 1$ [$g = 6/7$ and $g = 3/2$, respectively, for a nonrelativistic and relativistic degenerate gases of electrons (Ref. 41)], R is the star radius, $k = A/Z \approx (2 \text{ to } 2.5)$ is the ratio of the numbers of nucleons and protons in every nucleus, $N = M/AM_0$ is the number of nuclei in the volume of

a star, n_e is the concentration of degenerate electrons, G is the Newton gravitational constant, $M_0 = [M_p + M_n(k-1)]/k \approx (M_p + M_n)/2$ is the mean mass of a nucleon.

This energy corresponds to the negative (i.e., directed to the center) pressure on the star surface,

$$\begin{aligned} P_G &= -(dW_G/dV) = -gGM^2/4\pi R^4 \\ &= -gG(kM_0)^2 \left[(4\pi/3)^{1/3}/3 \right] (NZ)^{2/3} n_e^{4/3} \\ &= -gG(kM_0)^{4/3} \left[(4\pi/3)^{1/3}/3 \right] M^{2/3} n_e^{4/3}. \end{aligned} \quad (11.54)$$

It is seen from Eq. 11.54 that the adiabatic exponent for the gravitational pressure (the exponent characterizing the dependence of the pressure on the concentration of electrons) is $\gamma = \gamma_G = 4/3$.

The process of compression is hampered by the positive (directed from the center) kinetic (Fermi) pressure of the degenerate electron gas.

The total energy and pressure of a nonrelativistic degenerate gas of electrons are

$$W_{Fnr} = (3\pi^2)^{2/3} (3\hbar^2/10m_e) NZ n_e^{2/3}, \quad (11.55)$$

$$P_{Fnr} = (3\pi^2)^{2/3} (\hbar^2/5m_e) n_e^{5/3}. \quad (11.56)$$

Here, m_e is the electron mass.

The adiabatic exponents for a nonrelativistic gas of electrons are $\gamma = \gamma_{Fnr} = 5/3$ and $\gamma_G < \gamma_{Fnr}$.

Respectively, for a relativistic degenerate gas of electrons, we have

$$W_{Fr} = (3\pi^2)^{1/3} (3\hbar c/4) NZ n_e^{1/3}, \quad (11.57)$$

$$P_{Fr} = (3\pi^2)^{1/3} (\hbar c/4) n_e^{4/3}. \quad (11.58)$$

The adiabatic exponents for a relativistic gas of electrons are $\gamma = \gamma_{Fr} = 4/3$, $\gamma_G = \gamma_{Fr}$.

The boundary of the nonrelativistic and relativistic degenerate gases corresponds to the threshold concentration of degenerate electrons

$$n_{e0} \approx (m_e c / \hbar)^3 / 3\pi^2 \approx 10^{30} \text{ cm}^{-3}. \quad (11.59)$$

The comparison of quantities Eqs. 11.54 and 11.56 yields that, in the case of the insufficiently strong preliminary compression of a degenerate electron gas corresponding to a nonrelativistic degenerate electron gas, the state of a star is stabilized on the basis of the condition

$$P \equiv P_G + P_{Fnr} = 0 \quad (11.60)$$

at the electron density

$$n_{\text{est}} = (5gGm_e/9\pi\hbar^2)^3 4M^2 (kM_0)^4. \quad (11.61)$$

A basically other situation will be for a star preliminarily compressed to the state of a degenerate relativistic gas of electrons. By comparing quantities Eqs. 11.54 and 11.58, we find that the sign of the total pressure on the star surface,

$$P \equiv P_G + P_{Fr} = \left\{ (3\pi^2)^{1/3} (\hbar c/4) - gG(kM_0)^{4/3} \left[(4\pi/3)^{1/3}/3 \right] M^{2/3} \right\} n_e^{4/3}, \quad (11.62)$$

depends on the star mass, but does not depend on the concentration of the degenerate electron gas.

In particular, if the star mass M exceeds the critical value

$$M_{\text{cr}} = (9/16)(3\pi/g^3)^{1/2} (\hbar c/G)^{3/2} / (kM_0)^2 \approx 1.45M_\odot, \quad (11.63)$$

then the pressure related to the gravitational forces exceeds the kinetic pressure of the electron gas. Here, $M_\odot \approx (\hbar c/G)^{3/2} / 2M_0^2 \approx 2 \times 10^{33} \text{ g}$ is the sun's mass.

For this reason, the star surface is subjected to the action of the uncompensated negative pressure $P < 0$ increasing with the concentration of electrons. This result corresponds to the gravitational collapse of a star. If $M < M_{\text{cr}}$, the gravitational forces cannot resist to the kinetic pressure of the relativistic degenerate gas of electrons, and the expansion of the star occurs with the simultaneous decrease in the concentration of electrons to that corresponding to the state of a degenerate nonrelativistic gas. This process is stabilized at level Eq. 11.61.

In the frame of the comparative analysis of the gravitational and Coulomb collapses, we carry out an additional analysis of the evolution of an atom to a collapsing self-compressing Wigner–Seitz cell containing the nucleus and the degenerate gas of electrons upon an increase in the pressure. This consideration supplements naturally the general results derived above in Sec. 11.1.3

In the case of a neutral atom, the interaction of electrons with the nucleus depends on the way of the formation of a neutral atom.

For a free gas of electrons (not compressed by the external forces), the Coulomb interaction energy leads to the creation of ordinary atoms with fixed electron shells around every nucleus.

For a compressed atom, the result depends on the degree of the initial compression. Though these questions were considered in details in

Sec. 11.1.3, we will analyze those aspects which are related to the comparison of gravitational and Coulomb collapses.

If we take that the energy of the pairwise interaction of every electron with nuclei is defined by the standard formula for the Coulomb law

$$U_Q = -e\varphi(r) = -Ze^2/r, \quad (11.64)$$

the total interaction energy of the system of electrons with the nucleus in the case of a compressed quasiatom in the form of a Wigner–Seitz cell is mainly defined by two mechanisms:

(a) Coulomb attraction of electrons and the nucleus (Eq. 11.27),

$$\begin{aligned} W_{Q(en)} &= \int_0^R U_Q(r) n_e 4\pi r^2 dr = -(3/2)(Z^2 e^2/R) \\ &= -(3/2)(4\pi/3)^{1/3} Z^{5/3} e^2 n_e^{1/3}; \end{aligned} \quad (11.65)$$

(b) Coulomb mutual repulsion of electrons (Eq. 11.28),

$$\begin{aligned} W_{Q(ee)} &= \int_0^R Q(r) dQ(r)/r = (3/5)Z^2 e^2/R \\ &= (3/5)(4\pi/3)^{1/3} Z^{5/3} e^2 n_e^{1/3}. \end{aligned} \quad (11.66)$$

The contribution of the exchange energy to the process of interaction is insignificant and can be omitted in the first approximation.

The summary energy of the Coulomb interaction in the volume of a compressed quasi-atom is

$$W_Q = -(9/10)Z^2 e^2/R = -(9/10)(4\pi/3)^{1/3} Z^{5/3} e^2 n_e^{1/3}. \quad (11.67)$$

This energy corresponds to the negative (directed to the nucleus) “Coulomb” pressure on the surface of a compressed quasi-atom.

$$P_Q = -(3/10)(4\pi/3)^{1/3} Z^{2/3} e^2 n_e^{4/3}. \quad (11.68)$$

It is seen from Eq. 11.68 that the adiabatic exponent for the Coulomb pressure (like for the gravitational one in Eq. 11.54) $\gamma = 4/3$.

This pressure counteracts the kinetic pressure of the gas of nonrelativistic or relativistic electrons, Eq. 11.56 or 11.58, with the relevant adiabatic exponents $\gamma = 5/3$ and $\gamma = 4/3$.

In the case of nonrelativistic electrons, the total pressure on the surface of a Wigner–Seitz cell is determined upon the comparison of Eqs. 11.56 and 11.68 and has the form

$$P_{\Sigma} = \left[(3\pi^2)^{2/3} (\hbar^2/5m_e) n_e^{1/3} - (3/10)(4\pi/3)^{1/3} Z^{2/3} e^2 \right] n_e^{4/3}. \quad (11.69)$$

Equation 11.69 yields that, in the case of a nonrelativistic gas of degenerate electrons, there exists their optimum concentration

$$n_e = 4Z^2 (e^2 m_e / 2\pi \hbar^2)^3 \approx 10^{24} Z^2 \text{ cm}^{-3}, \quad (11.70)$$

at which $P_{\Sigma} = 0$. In this case, a compressed quasiatome is in the state of stable equilibrium. Because the maximum value of n_e for any known nuclei is significantly less than the threshold concentration $n_{e0} \approx 10^{30} \text{ cm}^{-3}$, the stable state for such nuclei corresponds to nonrelativistic electrons.

Such a situation is analogous to the stable state of a gravitational object, whose shell contains a nonrelativistic degenerate gas of electrons.

At a significantly greater concentration of degenerate electrons (i.e., in the case of a relativistic degenerate electron gas), the total pressure on the surface of a Wigner–Seitz cell can be found from Eqs. 11.58 and 11.68 as

$$P_{\Sigma} = \left[(3\pi^2)^{1/3} (\hbar c/4) - (3/10)(4\pi/3)^{1/3} Z^{2/3} e^2 \right] n_e^{4/3}. \quad (11.71)$$

It is seen from Eq. 11.71 that, in the case of the extremely compressed degenerate electron gas, there exists the threshold charge of a nucleus

$$Z_{\text{cr}} \approx (\pi/14)^{1/2} (5/2)^{3/2} (\hbar c/e^2)^{3/2} \approx 1.87 (\hbar c/e^2)^{3/2} \approx 3000, \quad (11.72)$$

beginning from which the summary pressure on the surface of every Wigner–Seitz cell becomes negative ($P_{\Sigma} < 0$), which leads to the irreversible compression of the electron-nucleus system.

Such a result corresponds to the threshold Coulomb collapse of the electron-nucleus plasma which is completely analogous to the threshold gravitational collapse. Moreover, the threshold charge of a nucleus $Z_{\text{cr}} \approx 3000$ is analogous to the critical mass of a star, $M_{\text{cr}} \approx 1.45 M_{\odot}$.

It seems at the first glance that the Coulomb collapse is basically unrealizable, because the nuclei with unrealistically large electric charge Z_{cr} are needed for its realization. For the nuclei with a lesser charge, the energy of the mutual Coulomb attraction of the nucleus and electrons is insufficient to overcome the repulsive force related to the kinetic pressure of the gas of electrons. The unique manifestation of the Coulomb interaction at $Z \ll Z_{\text{cr}}$ can be a small diminution of the total pressure of the degenerate relativistic

electron gas in a compressed quasiatome, being at most 2% to 3%. Such a conclusion was made on the basis of the use of the law of pairwise interaction Eq. 11.64 in works in Refs. 42, 43, 44, 45 devoted to the consideration of the electron–nuclear interaction in the problem of stability of astrophysical objects.

In fact, as shown by the comprehensive analysis performed in Sec. 11.1.3, this conclusion is incorrect, and the process of Coulomb collapse is possible also for nuclei with small charge $Z \ll Z_{\text{cr}}$ (including the “ordinary” nuclei with $Z \leq 92$)!

The point is in that, to obtain the basic relations Eqs. 11.67, 11.68, 11.71, and 11.72, we use the formula for the energy of the pairwise Coulomb interaction of the nucleus with an electron Eq. 11.64 which is true only for immovable (or slowly moving) electrons and only in the case where this energy is significantly less than the rest energy of electrons, $m_e c^2$. It is obvious that both these conditions are not satisfied in the zone of action of the gravitational collapse.

As follows from the analysis of the Dirac equation given in Sec. 11.1.3, the full expression for the potential energy of an electron in the field of the potential $\varphi(r) = Ze/r$ significantly differs from the standard Coulomb law Eq. 11.64 and has the form Eq. 11.12

$$U_{Q\text{eff}}(r) \equiv U_{Q\text{eff}}(r, E) = -e\varphi(r)(E/m_e c^2) - [e\varphi(r)]^2/2m_e c^2. \quad (11.73)$$

The importance of such a circumstance can be estimated by comparing the usually used “linear” formula Eq. 11.64 for the potential energy of an electron $U_Q(R)$ and the nonlinear part $U_{Q\text{eff}}^{NL}(R) = -(e\varphi(r))^2/2m_e c^2$ of the relation for $U_{Q\text{eff}}(R)$ near the nucleus surface (at $R = 5 \times 10^{-13}$ cm). After the direct substitution, the parameters are as follows:

$|U_Q(R)| = 0.28Z \text{ MeV} = 7.3 \text{ MeV}$, $|U_{Q\text{eff}}^{NL}(R)| = 0.08Z^2 \text{ MeV} = 54 \text{ MeV}$ at $Z = 26$ and $|U_Q(R)| = 25.7 \text{ MeV}$, $|U_{Q\text{eff}}^{NL}(R)| = 680 \text{ MeV}$ at $Z = 92$.

It is seen that the use of the modified Coulomb law leads to a very considerable (by several orders) increase in the energy of the Coulomb interaction of the nucleus and electrons. The same effect induces a very significant change in the Coulomb pressure on the surface of a compressed quasiatome as compared to the estimations derived on the basis of Eq. 11.64. Analogous variations in $U_{Q\text{eff}}(r)$ occur upon the account of relativistic effects.

These peculiarities of the Coulomb interaction in the volume of a compressed relativistic degenerate gas were not considered in the previous works and allow us to predict the existence of a Coulomb collapse in a preliminarily compressed electron-nucleus plasma.

The account of nonlinear and relativistic effects leads to a sharp increase in the adiabatic exponent defining the increase in the compressing

Coulomb pressure with the density: from the initial value $\gamma_Q = 4/3$ which is too small to overcome the kinetic pressure $P_{F(nr)} \sim n_e^{\gamma_F}$ of a nonrelativistic ($\gamma_F = 5/3$) electron gas, as well as the pressure of the relativistic gas ($\gamma_F = 4/3$) for nuclei with $Z \ll 3000$, to the final value $\gamma_Q = 5/3$. Such a change in the adiabatic exponent results in that, at a definite threshold concentration of electrons in the degenerate relativistic gas of electrons, the Coulomb pressure exceeds the kinetic pressure and ensures the possibility to realize the Coulomb collapse not only for the nuclei with anomalously large charge $Z > 3000$, but also for the ordinary nuclei with $Z < 92$.

It is worth noting that just these two circumstances (the influence of the relativistic motion and a nonlinear character of the Coulomb interaction of particles) basically distinguish the Coulomb collapse from the gravitational one: the gravitational collapse of a star occurs at the expense of the gravitational interaction of nonrelativistic nuclei and is characterized by the usual “linear” classical formula for the potential energy of the Newton interaction $U_g = -GMm_e/r$, whereas the Coulomb collapse occurs at the expense of the electromagnetic interaction of relativistic electrons in the region of manifestations of the nonlinear Coulomb interaction $U_{Q\text{eff}}(r)$.

Thus, the conditions for a realization of the effect of Coulomb collapse can be satisfied in the process of preliminary compression of the shell of a star realizing in the course of the gravitational collapse.

Cosmological Consequences of the Possible Coulomb Collapse in a Gravitating Star. The derived estimations yield that, in order to initiate the Coulomb collapse in the volume of every neutral Wigner–Seitz cell, the preliminary attainment of a high density of the degenerate electron gas is needed.

Such a density can be reached with excess in the process of gravitational collapse preceding to the formation of a neutron star or a black hole (as known, the density of a degenerate electron gas can reach 10^{34} cm^{-3} on the inner surface of the shell of a heavy gravitating star). It follows from the above-performed estimations that, in the shell of a star being in the state of gravitational collapse, the Coulomb collapse is possible for various nuclei (including the relatively light nuclei with $Z = 26$ to 14).

Due to the last circumstance, the influence of the Coulomb collapse on the evolution of a heavy star can be essential. Is the Coulomb collapse possible in other astrophysical objects?

We have shown above that the critical value of the mean kinetic energy of electrons in the degenerate gas for the heaviest “ordinary” nuclei (e.g., nuclei U or Pb which are always present in small amounts in the stellar atmosphere) corresponds to the value $U_{e\Sigma}/Z \approx$

0.65 to 0.8 MeV/electron and requires a relatively small starting density $n_{e(\text{eff})} \approx [(64/81\pi^2)/(\hbar c)^3] U_{e\Sigma}/Z \approx (1 \text{ to } 2) \times 10^{30} \text{ cm}^{-3}$ which increases up to $n_{e(\text{crf})} \approx (2 \text{ to } 3) \times 10^{32} \text{ cm}^{-3}$ in a Wigner–Seitz cell. Because such a density $n_e \approx 10^{30}$ to $2 \times 10^{30} \text{ cm}^{-3}$ is quite attainable in stable white dwarfs, the Coulomb collapse is also possible in the last objects for the heaviest nuclei. It is natural that, due to a very small concentration of heavy nuclei, the influence of this effect on the state of white dwarfs is slight and is not comparable with the influence on a heavy star, being in the state of gravitational collapse.

Besides the natural tandem of the gravitational and Coulomb collapses of a heavy star, where the former stimulates the start of the latter, the condition for a realization of the Coulomb collapse in stars can be also related to the process of formation of a spherical wave (a thin spherical layer) propagating to the center of a star with anomalously high density of the degenerate plasma and can run in the absence of a gravitational collapse. During the accelerated movement to the center, this wave evolves rapidly to the state of propagating Coulomb collapse, which stimulates naturally a great variety of nuclear reactions. A mechanism of formation and acceleration of such a wave will be thoroughly considered in the next chapter.

The above-considered process of formation of the electron-nucleus Coulomb collapse, the starting conditions for which are satisfied at the expense of the compression of the electron-nucleus plasma in the preceding gravitational collapse of a massive star or at the expense of a specific self-compression in the volume of a thin spherical wave, allows us to advance a mechanism of nucleosynthesis alternative to the traditional models. Since the evolution of the electron-nucleus Coulomb collapse is accompanied by a sharp displacement of the region of energy-gained synthesis toward superheavy nuclei with $A \gg 60$ and the screening of a nucleus with decrease in the height of a Coulomb barrier, such nuclei will be inevitably formed in the compressed shell of a collapsing star.

In the process of fast growth of the mass of a superheavy nucleus (at the expense of the absorption of light nuclei from the environment), the great binding energy is released. A part of this energy can leave the nucleus in the form of γ -emission and neutrinos. The other part can stimulate the emission of various “ordinary” nuclei with $A < 200$ to 230 by a growing superheavy nucleus, by forming the channel of synthesis combined with accompanying fission (Fig. 11.9). In the course of the realization of such “fusion-fission” chain, a great variety of nuclei can be formed.

It is obvious that the highest probability is characteristic of the emission of such maximally stable (at a specific density of the compressed relativistic electron gas) nuclei which are created in the form of clusters through

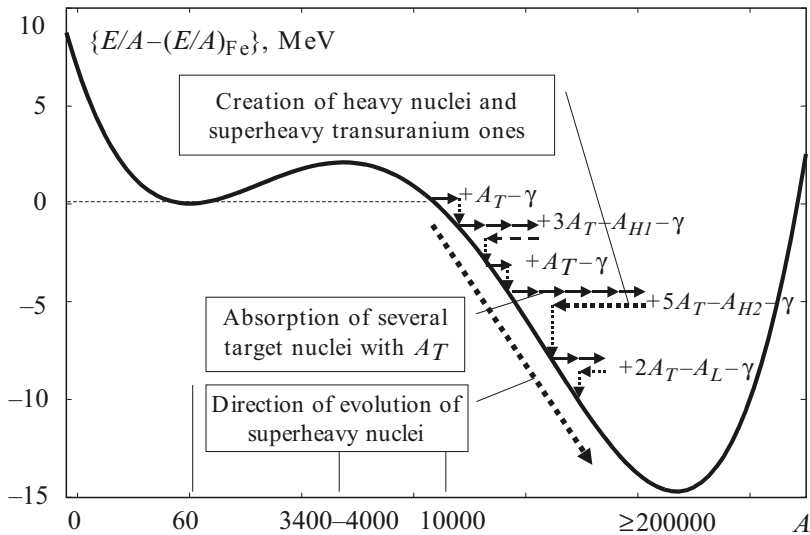


Fig. 11.9. Evolution of superheavy nucleus—absorption of target nuclei and creation of different nuclei (from H up to stable transuranium nuclei).

fluctuations in the volume of a growing superheavy nucleus. If such a cluster is insufficiently stable and has a moderate binding energy, it will rapidly decay in the volume of a very big superheavy nucleus at once after the creation. This question was considered above in more details. Only stable clusters will be emitted from the nucleus, which corresponds to the formation of nuclei close to the stability line. Such a process testifies to that every growing superheavy nucleus can be, in essence, a microreactor which can absorb light nuclei (e.g., hydrogen) and synthesize others (as a rule, stable ones).

The channel of emission of γ -quanta, which is not related to the ejection of relatively light nuclei from a growing superheavy nucleus, is alternative to the channel of synthesis of nuclei upon the absorption of nuclei of a target. For this reason, the mass of the superheavy nucleus will be gradually increased.

Because the considered process of synthesis of nuclei of any mass is not connected with the chain of subsequent transformations (like that in the s - and r -processes), the problems of synthesis of neutron-deficient nuclei and nuclei of the elements Li, Be, and B do not generally arise. These nuclei are created in the natural way according to the considered “fusion-fission” mechanism.

Of course, such a scenario of nucleosynthesis must be thoroughly compared with the law of abundance of isotopes and elements (especially in the part characterizing the abundance of even-even and other stable nuclei).

However, already the cursory analysis confirms that the abundance of elements correlates well with their stability and, hence, with the mechanism of their creation. It is natural that, in addition to such a “fusion–fission” mechanism, a very significant role in the creation of elements is played by the s - and r -processes, which also influences the abundance of elements and isotopes.

The derived results allow us to formulate a conjectural answer to the question about the mechanism of the local release of a very great energy $E_S \approx 10^{51}$ erg necessary for the separation of the shell of a star in the process of gravitational collapse and its acceleration to a very high velocity $v_S \approx 10^9$ cm/s.

As mentioned above, neither a shock wave, nor a flux of neutrinos from the center of a gravitational collapse can ensure the transfer of such great energy from the central region of the gravitational collapse to the inner surface of the shell. If we take into account the possibility of a realization of the sequential tandem of the gravitational and Coulomb collapses, the question about the source of this energy can be solved in a sufficiently cardinal way: the energy E_S is not transferred in some way from the center of the gravitational collapse, but it is a result of numerous acts of the Coulomb collapse in the volume adjacent to the inner surface of the shell.

In such a model, the process of Coulomb collapse begins in that part of the shell of a gravitating star, where the density of the degenerate relativistic electron gas reaches the threshold value Eq. 11.40 $n_e \approx 3 \times 10^{36} Z^{-2} \text{ cm}^{-3}$. In particular, on the outer boundary of the iron core of a gravitating star, the threshold condition for the Coulomb collapse corresponds to the density of electrons

$$n_{e(\text{Fe})} \approx 4.4 \times 10^{33} \text{ cm}^{-3}. \quad (11.74)$$

It is expedient to emphasize that the Coulomb collapse can start at a considerably lesser density and can be formed in the process of propagation of a spherical shock wave with the anomalously high density of electrons which converges to the center. This process will be comprehensively considered in the following chapter.

It is easy to verify that the energy released upon the Coulomb collapse is sufficient for the separation of the shell of a star and its subsequent acceleration to a high velocity.

As shown above, in the process of “falling of electrons onto the nucleus” associated with the Coulomb collapse, the energy reaching c is released. Because this energy is a result of the Coulomb interaction of electrons and the nucleus, it is released, most likely, in the form of hard X-ray

emission with a maximum at

$$\hbar\omega \approx \hbar c \langle 1/r \rangle = (9\pi/2)^{1/3} \hbar c (n_{\text{ecr}}/Z)^{1/3} \approx 6/Z^{1/3} \text{ MeV}. \quad (11.75)$$

The path length for such an emission in the very dense medium of the shell of a star does not exceed tenths of centimeter, which leads to a local heating of this part of the shell.

We will perform simple estimations of the energy released in such a process occurring in a gravitating star subjected to the gravitational collapse.

In any case, the mass of the core of such a star always exceeds the Chandrasekhar threshold equal to $1.45M_{\odot}$. The total mass of a gravitating star corresponds usually to the interval (6 to $10M_{\odot}$).

If the reactions of the Coulomb collapse occurring in the shell of a star involve only 0.1% the total mass of a star ($\Delta M \approx 0.01M_{\odot} = 2 \times 10^{31} \text{ g}$), then the total released energy is equal to

$$\Delta W \Delta M / k M_n \approx 10^{51} \text{ erg}. \quad (11.76)$$

This energy is quite sufficient for the heating and separation of the shell of a collapsing star and for its subsequent acceleration up to the velocity $v_s \approx 10^9 \text{ cm/s}$.

Thus, the mechanism based on the sequential tandem of the gravitational and Coulomb collapses of a part of the shell of a heavy star allows us to answer many unsolved problems of astrophysics and nucleosynthesis.

In conclusion, we note that the question about the limiting degree of self-compression of the system of degenerate relativistic electrons in a compressed quasiautom (i.e., about the attainment of the maximum ratio n_e/n_p , the possibility to weaken the interaction of protons in a nucleus, and the limiting mass of a nucleus can be solved only within the full analysis of the whole variety of electron-nucleus transformations (including the processes of protonization and neutronization and formation of electron-positron pairs) for specific nuclei. Some of these questions will be considered in the following section. Undoubtedly, a certain role should be played by the process of synthesis of giant nuclei with mass numbers $A \gg 1000$ which was considered by A. Migdal with regard to the mechanism of pion condensation for such nuclei (see the discussion in the above sections). It is obvious that since such a mechanism can be realized, according to A. Migdal, only for the nuclei with $Z > 1600$, the above-considered scenario of the synthesis of superheavy nuclei by means of the Coulomb collapse on the basis of "ordinary" nuclei has no real alternative in the region $100 < Z < 1600$.

11.1.6. Basic Reactions in the Collapse of the Electron-Nucleus System

Threshold Conditions for the Creation of Electron-Positron Pairs in the Collapse Zone. Of great interest is the question about the limiting value of the binding energy of the compressed medium under consideration representing the system of Wigner-Seitz cells containing the degenerate electron gas. As shown above (see, e.g., Figs. 11.4 and 11.5), the interaction of nuclei and the extremely compressed degenerate relativistic electron gas induces the effect of “falling of electrons onto the nucleus” and the collapse of the electron-nucleus system, which leads to the release of a huge amount of binding energy. This energy can be released in various channels.

In nuclear systems, one of the main channels of the energy release is connected with the instant hard γ -emission with the upper energy limit corresponding to the value of the released binding energy. The alternative channel of the energy release can be connected with the creation of electron-positron pairs and with the basic reactions of neutronization and protonization of a nucleus located in the volume of a Wigner-Seitz cell.

Consider the possibility of a realization and the conditions for the running of each of these processes.

Upon the release of a relatively moderate binding energy, the main mechanism of the removal of energy from the volume of a compressed system should involve the instant γ -emission. However, in an axially compressed symmetric Wigner-Seitz cell, such an electromagnetic emission corresponds to monopole electromagnetic transitions, whose probability is very small. Such transitions correspond, in particular, to nonemitting axially pulsating spherical charges. The emission from such a system can occur only at the expense of a slight difference of the model spherical Wigner-Seitz cell from a real system.

It seems at the first glance that an increase in the binding energy of a cell, which consists of the degenerate gas of relativistic electrons and the nucleus and undergoes the compression, must lead to the stage-by-stage creation of a great number of electron-positron pairs. On the other hand, it is obvious that an increase in the Fermi energy of the degenerate electron gas in this Wigner-Seitz cell can lead to a partial suppression of this process under its compression. Consider the process of creation of pairs in more details and with regard to the above-mentioned factors.

We recall that the creation of a pair will be possible only if the energy of the electron from this pair is sufficient for the population of the lowest level from the free quantum ones. This condition yields that the minimum energy of a pair upon its creation in the volume of the degenerate electron

gas must exceed the doubled Fermi energy of this gas,

$$E_{ee} \equiv 2(m_e^2 c^4 + p_F^2 c^2)^{1/2} \geq 2E_F. \quad (11.77)$$

The threshold of the creation of a pair in such a system is determined from the condition of the equality of the binding energy of one electron, $U_{e\Sigma}/Z$, and $2E_F$ in the volume of the superdense degenerate relativistic gas.

Here, p_F is the limiting (Fermi) momentum of an electron in the volume of a Wigner–Seitz cell. In the considered range of densities of the superrelativistic gas of electrons, the Fermi energy is equal to $E_F = p_F c = (3\pi^2)^{1/3} \hbar c n_e^{1/3}$.

In Fig. 11.10, we present the plots of the energy of a degenerate electron gas $U_{e\Sigma}/Z$ in the volume of a compressed quasiautom calculated by formula Eq. 11.35 at the extremely high density of this gas and the threshold energy of creation of an electron–positron pair, $2|E_F(n_e)|$, at the same region of values of n_e . It is seen from this figure that the increase in the electron gas density in the scope of the entire region of values, $n_e < 2 \times 10^{35} \text{ cm}^{-3}$, leads to a significant increase in the binding energy $U_{e\Sigma}/Z$, but does not cause the automatic attainment of the threshold of creation of an electron–positron pair. This is conditioned by the fact that, at the same values of n_e , the threshold energy of creation of a pair also increases synchronously with the binding energy (though by the weaker law $2E_F \sim n_e^{1/3}$).

The lowest threshold of creation of a pair is determined from the condition $U_{e\Sigma}(n_e)/Z_{\text{max}} = 2|E_F(n_e)|$ and corresponds to a Wigner–Seitz cell, at the center of which a nucleus with maximum charge is located. From stable nuclei, we indicate uranium with $Z = 92$. For such a cell, the threshold specific and total binding energy and the threshold electron density, respectively, are $(U_{e\Sigma}/Z)_{Z=92(ee)} \approx 77 \text{ MeV}$, $(U_{e\Sigma})_{Z=92(ee)} \approx 7 \text{ GeV}$, $(n_e)_{Z=92(ee)} \approx 2 \times 10^{35} \text{ cm}^{-3}$.

With decrease in the nucleus charge Z , the threshold of creation of a pair significantly increases.

For Pb, we get $(U_{e\Sigma}/Z)_{Z=82(ee)} \approx 81 \text{ MeV}$, $(U_{e\Sigma})_{Z=82(ee)} \approx 6.7 \text{ GeV}$, $(n_e)_{Z=82(ee)} \approx 3 \times 10^{35} \text{ cm}^{-3}$;
for Yb with $Z = 70$, $(U_{e\Sigma}/Z)_{Z=70(ee)} \approx 88 \text{ MeV}$, $(U_{e\Sigma})_{Z=70(ee)} \approx 6.1 \text{ GeV}$, $(n_e)_{Z=70(ee)} \approx 4 \times 10^{35} \text{ cm}^{-3}$;
for Sn with $Z = 50$, $(U_{e\Sigma}/Z)_{Z=50(ee)} \approx 100 \text{ MeV}$, $(U_{e\Sigma})_{Z=50(ee)} \approx 5 \text{ GeV}$, $(n_e)_{Z=50(ee)} \approx 8 \times 10^{35} \text{ cm}^{-3}$;
for Fe, $(U_{e\Sigma}/Z)_{Z=26(ee)} \approx 150 \text{ MeV}$, $(U_{e\Sigma})_{Z=26(ee)} \approx 3.9 \text{ GeV}$, $(n_e)_{Z=26(ee)} \approx 2 \times 10^{36} \text{ cm}^{-3}$;
and, for Si, $(U_{e\Sigma}/Z)_{Z=14(ee)} \approx 215 \text{ MeV}$, $(U_{e\Sigma})_{Z=14(ee)} \approx 3 \text{ GeV}$, $(n_e)_{Z=14(ee)} \approx 10^{37} \text{ cm}^{-3}$.

Upon the attainment of such a threshold, the superrelativistic electrons and positrons with energies exceeding the Fermi energy will leave the region of collapse, which corresponds to values 38 to 107 MeV. We note that these values of the energy define, in essence, the minimum energy of electrons and positrons upon their creation in the composition of pairs. But if the process of creation of pairs will “delay” in the course of time from the process of collapsing self-compression, then the energy of the created electrons and positrons can be considerably greater and can reach several GeV.

In addition, these results testify to that the contribution of the Coulomb energy of the interaction of the superrelativistic gas of electrons and the nucleus to the total energy of every Wigner–Seitz cell is quite comparable to analogous contributions to this energy from the side of the processes defining the surface $\varepsilon_2/A^{1/3}$, Coulomb $\varepsilon_3 Z^2/A^{4/3}$, and symmetry-related $\varepsilon_4(A/2 - Z)^2/A^2$ energies of the nucleus. Moreover, the contribution of the Coulomb energy of such an electron–proton interaction can become dominant in a number of cases.

We note that though the mean distance between protons and electrons $\langle r_{pe} \rangle$ is considerably more than the mean distance $\langle r_{pp} \rangle$ of protons in a nucleus, the comparability of the electron–proton interaction with

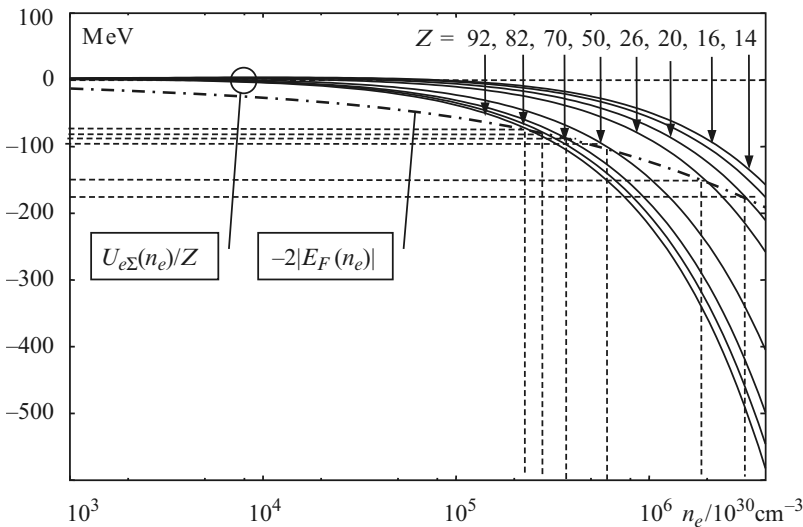


Fig. 11.10. Comparison of the energy of a degenerate electron gas $U_{e\Sigma}/Z$ in the volume of a compressed quasi-atom and the energy of creation of an electron–positron pair $2E_F(n_e)$ in the region with an extremely high density of the degenerate electron gas.

intranuclear interactions does not lead to a paradox. The point is in that the though protons in a nucleus are close each to other ($r_{pp} \ll r_{pe}$), they have a greater mass, and their velocity in the volume of a nucleus is small relative to the light velocity. As a result, the Coulomb interaction of protons is always characterized by the standard “linear” potential $\varphi(r_{pp})$. At the same time, the electrons in the volume of a compressed Wigner–Seitz cell are relativistic particles and, in addition, have a considerably lesser rest mass, which leads, at once, to that their interaction is characterized by the nonlinear [relative to the potential $\varphi(r_{pe})$] law Eq. 11.13.

If the electron gas density is less than the critical value $n_{eZ} \approx 2 \times 10^{35} \text{ cm}^{-3}$ and if this leads to the prohibition of the creation of pairs, then the main mechanism of the release of energy from the volume of a compressed system should involve the instant γ -emission. As noted above, such an emission will be significantly suppressed due to the symmetry of a compressed Wigner–Seitz cell.

Under these conditions (with the suppressed γ -emission and the Fermi prohibition of the creation of pairs), the released binding energy can be almost completely transferred to intranuclear processes and has the essential effect on the synthesis and transmutation of nuclei. In addition, such “nonutilized” store of a great binding energy (the energy not released by the channels of γ -decay and the creation of electron–positron pairs in the scope of every Wigner–Seitz cell) reaching several GeV per one atom of the type of Cu can become one of the reasons for the acceleration of the spherical plasma layer (shell). This shell can be formed in a macroscopic target and accelerated upon the action of a symmetric coherent driver on it. This process will be considered in the following chapter, and its realization is one of the premises for the attainment of the threshold density of a compressed electron gas, which induces the realization of the process of self-controlled Coulomb collapse.

Neutronization and Protonization of Nuclei in the Collapse Zone.

The other important process, which can run in the volume of a compressed Wigner–Seitz cell with a great density of degenerate electrons, is the transformation of nuclei at the expense of the direct weak interaction electrons with nucleons. We note at once that, by virtue of a number of reasons, the question about the directivity of such a process (neutronization or protonization) requires a deep analysis, because its directivity depends on several mutually excluding factors and presents a distinctive dilemma; see Ref. 39.

On the one hand, starting from the traditional approach frequently used in nuclear astrophysics (e.g., Ref. 44), we should recognize that the

total energy of the system is significantly decreased by the process of neutronization



This process occurs in the volume of the relativistic degenerate gas with a great Fermi energy with the transformation of one of the protons of the nucleus into a neutron (the process runs with energy consumption) and with the simultaneous absorption of one of the electrons from the Fermi-gas (in this case, the energy numerically equal to the Fermi energy is released).

On the other hand, it is necessary to note that the same process of neutronization decreases the nucleus charge Z and the number of electrons in a Wigner–Seitz cell, which causes the essential decrease in the energy of the “nonlinear” connection of the electron subsystem and the nucleus. Such a process leads obviously to an increase in the total energy of the same system.

If the first mechanism of the energy release is stronger than the second one, then the reaction of neutronization Eq. 11.78 will be of high priority.

In the opposite case (with the leading role of effects of the nonlinear Coulomb interaction), the main process is the protonization of nuclei,



Here, the spontaneous decay of a neutron with the creation of an additional proton is accompanied by the creation of a fast electron with energy exceeding the Fermi energy. The interaction energy of this “superfluous” electron with the “superfluous” nucleus charge will compensate the increase in the Fermi energy with excess.

To calculate the directivity of the reaction, we use Eq. 11.45 for the total energy of a compressed quasiatome $\Delta E_{en}/A$ normed per one nucleon. As a part, this formula includes the total energy of the degenerate electron gas $U_{e\Sigma}$ Eq. 11.35 composed from the total kinetic (Fermi) energy of the degenerate electron gas and the rest energy of individual electrons, $Zm_e c^2$.

As the standard value, we use the dependence of the total energy of a compressed quasiatome $\Delta E_{en}/A$ on the quantities A , Z , and n_e taken on the classical line of stability.

In Fig. 11.11, we present the total energy of a compressed quasiatome $\Delta E_{en}(A, Z, n_e)/A$ versus the atomic number A at a “low” density ($n_e \ll 3 \times 10^{30} \text{ cm}^{-3}$) and two values of a “moderately high” ($n_e = 10^{30} \text{ cm}^{-3}$ and $n_e = 3 \times 10^{30} \text{ cm}^{-3}$) density of the degenerate relativistic electron gas. The calculations correspond to the classical line of stability Eq. 11.46 for a nucleus with charge $Z = Z_{\text{opt}}(A)$ and to the case where the nucleus charge Z differs from $Z_{\text{opt}}(A)$ by $\Delta Z = 1$ toward the increase or decrease in Z .

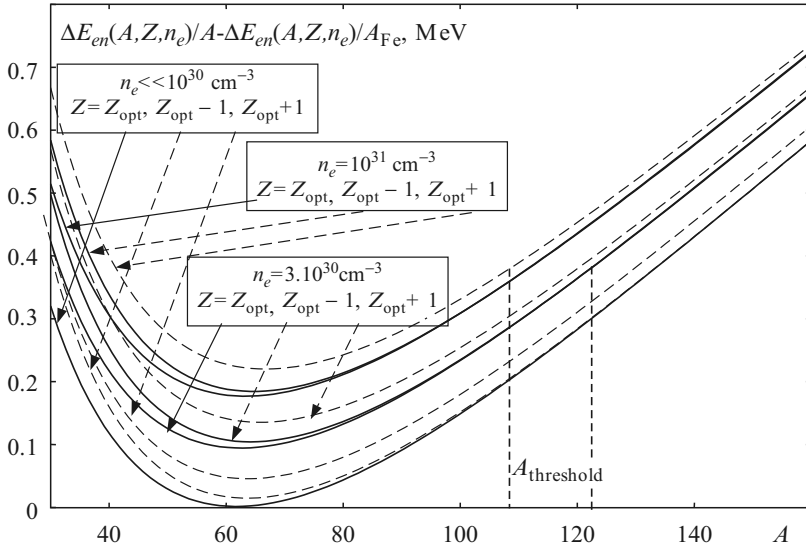


Fig. 11.11. Total energy of a compressed quasiautomon $\Delta E_{en}(A, Z, n_e)/A$ versus atomic number A at a relatively low density of degenerate electrons n_e on the line of stability Z_{opt} and upon a deviation from it (at $Z = Z_{opt} \pm 1$). Dashed vertical lines correspond to the threshold of neutronization of nuclei.

It is seen from the results presented in Fig. 11.11 that nuclei possess the maximum stability on the line of stability at a low density of the degenerate electron gas. This expected result is a distinctive test to verify the correctness of the used models and directly follows from that the plot characterizing the total energy of a system on the line of stability $\Delta E_{en}(A, Z_{opt}, n_e)/A$ for any A and at $n_e \ll 10^{30} \text{ cm}^{-3}$ is always lower than the analogous plot defining the total energy in the case of a deviation by ΔZ with any sign from the optimum charge $Z_{opt}(A)$.

With increase in the density n_e , the situation is abruptly changed. It follows from Fig. 11.11 that the “standard” line of stability Z_{opt} ceases to correspond to maximally stable nuclei with increase in the density of the degenerate electron gas. In this case, the directivity of the process of nuclear transmutation depends on the ratio of A for the initial nucleus and the threshold mass number $A_{threshold}$ characterizing the intersection point of the plots of the functions $\Delta E_{en}(A_{threshold}, Z_{opt}, n_e)/A_{threshold}$ and $\Delta E_{en}(A_{threshold}, Z_{opt} + \Delta Z, n_e)/A_{threshold}$. The values of $A_{threshold}$ depend on n_e and are presented in Fig. 11.11 by the vertical dashed lines at $\Delta Z = -1$.

For nuclei with a small mass ($A < A_{threshold}$), the plot of the total energy of the system $\Delta E_{en}(A, Z_{opt} + \Delta Z, n_e)/A$ as a function of the mass number A in the presence of a deviation ΔZ of any sign from Z_{opt} is

higher than the plot of the total energy on the line of stability Z_{opt} . This corresponds to that both the protonization and neutronization of nuclei become the unfavorable processes at $A < A_{\text{threshold}}$. Upon the use of nuclei with higher masses ($A > A_{\text{threshold}}$), the plot of the total energy of a Wigner–Seitz cell $\Delta E_{\text{en}}(A, Z_{\text{opt}} + \Delta Z, n_e)/A$ at $\Delta Z < 0$ lies lower than the plot of the total energy corresponding to the line of stability Z_{opt} . This means that, with increase in the density n_e and at $A > A_{\text{threshold}}$, the decrease in the nucleus charge in the volume of a neutral Wigner–Seitz cell is an energy-gained process, which corresponds to the neutronization of a nucleus. This result well agrees with the standard ideas of the directivity of β -processes in a relativistic degenerate electron gas.

A realization of the process of neutronization in the frame of the above-considered dilemma corresponds to that, at a moderately high density of the relativistic degenerate gas, a decrease in the Fermi energy upon the absorption of an electron by a proton is a more favorable process, than a decrease in the binding energy of electrons with the nucleus in the scope of a Wigner–Seitz cell upon such a change in the charge.

At the same time, this result corresponds to a certain interval of values of the electron gas density. If the density exceeds the upper boundary of this interval, the process of protonization becomes favorable.

In Fig. 11.12, we presented the total energy of a compressed quasihatom $\Delta E_{\text{en}}(A, Z, n_e)/A$ versus the atomic number A at a “low” density ($n_e \ll 3 \times 10^{30} \text{ cm}^{-3}$) and two values of a “very high” ($n_e = 10^{34} \text{ cm}^{-3}$ and $n_e = 1.5 \times 10^{34} \text{ cm}^{-3}$) density of the degenerate relativistic electron gas. The results of calculations correspond to both the classical line of stability Eq. 11.46 for a nucleus with charge $Z = Z_{\text{opt}}(A)$ and to the case where the nucleus charge Z differs from $Z_{\text{opt}}(A)$ by $\Delta Z = \pm 3$.

It is seen from Fig. 11.12 that, at a very high density n_e and upon the use of nuclei with large masses ($A > A_{\text{threshold}}$), the plot of the total energy of a Wigner–Seitz cell $\Delta E_{\text{en}}(A, Z_{\text{opt}} + \Delta Z, n_e)/A$ at $\Delta Z > 0$ lies below the plot of the total energy corresponding to the line of stability Z_{opt} . This means that, in this case at $A > A_{\text{threshold}}$, the increase in the nucleus charge (and, respectively, the same increase in the number of electrons) in the volume of a neutral Wigner–Seitz cell is an energy-gained process, which corresponds to the process of protonization of a nucleus.

It also follows from the same figure that the threshold value of the mass number $A_{\text{threshold}}$ decreases with increase in the density n_e (at $n_e = 10^{34} \text{ cm}^{-3}$, we have $A_{\text{threshold}} \approx 90$; and it decreases to $A_{\text{threshold}} \approx 80$ at $n_e = 1.5 \times 10^{34} \text{ cm}^{-3}$).

To confirm this tendency (a decrease in $A_{\text{threshold}}$ with increase in n_e), we study the process of nuclear transformations in a wider range of variations in the density of electrons.

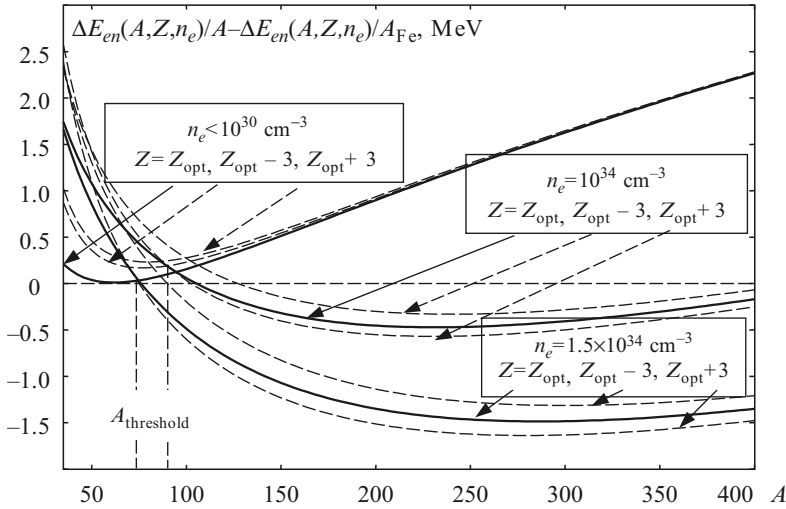


Fig. 11.12. Total energy of a compressed quasiautomon $\Delta E_{en}(A, Z, n_e)/A$ versus the atomic number A at a very high density of degenerate electrons n_e on the line of stability Z_{opt} and upon the deviation from it (at $Z = Z_{opt} \pm 3$).

In Fig. 11.13, we demonstrate the result of calculations and compare the energy of a neutral compressed Wigner–Seitz cell on the line of stability $\Delta E_{en}(A, Z_{opt}, n_e)/A$ (this is the standard result) and the energy of the same system $\Delta E_{en}(A, Z_{opt} + \Delta Z, n_e)/A$ in the case where its charge is increased by $\Delta Z = 1$ relative to Z_{opt} (i.e., under the protonization of a nucleus). The calculations were performed for the electron density changing in the wide limits: from $n_e = 3 \times 10^{34} \text{ cm}^{-3}$ to $n_e = 10^{36} \text{ cm}^{-3}$. The derived results allow us to determine the displacement of the protonization threshold $A_{threshold}$ upon a change of n_e . It is seen that the protonization threshold continuously decreases with increase in the density n_e : from $A_{threshold} \approx 75$ at $n_e = 3 \times 10^{34} \text{ cm}^{-3}$ corresponding to the chemical elements which are medium by weight to $A_{threshold} \approx 7$ at $n_e = 10^{36} \text{ cm}^{-3}$, which corresponds to lightest elements.

To illustrate the threshold character of the process of neutronization and the existence of the upper boundary for this process, Fig. 11.14 shows the dependence of the energy $\Delta E_{en}(A, Z, n_e)/A$ of a compressed quasiautomon on the atomic number A and the density of the gas of degenerate electrons n_e on the line of stability Z_{opt} and upon the deviation from it by $\Delta Z = \pm 1$.

These results allow us to determine the regularity of the process of neutronization as a function of the density of electrons. At the beginning (at a small excess of the density above the threshold value, $n_e \approx 10^{30} \text{ cm}^{-3}$),

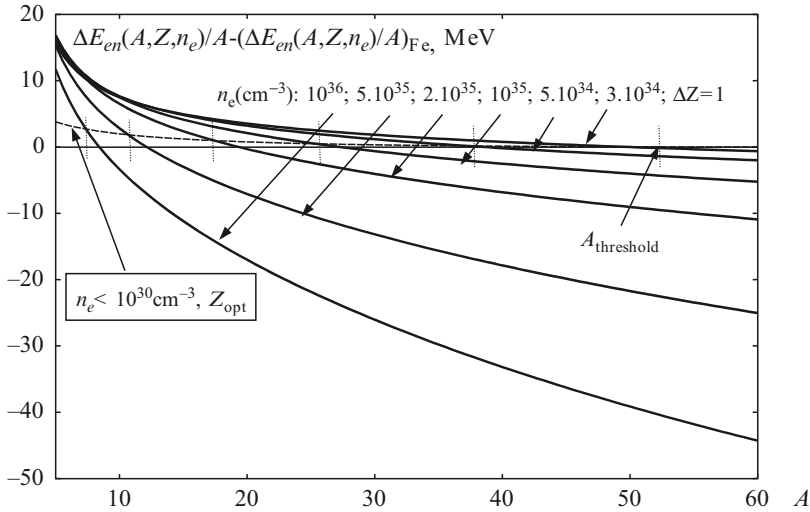


Fig. 11.13. Total energy $\Delta E_{en}(A, Z, n_e)/A$ of a compressed quasiatom versus atomic number A for various low densities of the gas of degenerate electrons n_e on the line of stability Z_{opt} upon the deviation from it by $\Delta Z = 1$. The first threshold of protonization ($\Delta Z = 1$) corresponds to the intersection of the curves $\Delta E_{en}(A, Z_{opt}, n_e < 10^{30} \text{ cm}^{-3})/A$ and $\Delta E_{en}(A, Z_{opt} + 1, n_e)/A$.

only the heaviest nuclei can participate in the process of neutronization. With increase in n_e , the region of neutronization expands onto lighter nuclei, but heavier nuclei begin to leave this region. This is conditioned by that the process becomes unfavorable by energy for the heavier nuclei due to a sharp increase in the Coulomb interaction with electrons. At last, upon the attainment of the upper limit of the “neutronization corridor”, $10^{30} \text{ cm}^{-3} \leq n_e \leq 10^{35} \text{ cm}^{-3}$, the neutronization becomes unfavorable for light nuclei as well. At such a value of n_e , a decrease in the Fermi energy turns out to be insufficient for the compensation of a decrease in the binding energy of electrons with the nucleus due to a weakening of the Coulomb interaction and for the compensation of the energy expenses on the increase in the nucleon mass (a neutron instead of a proton).

Thus, the process of neutronization can be the basic process of transformations of nuclei in the degenerate electron gas with the density corresponding to the “neutronization corridor”, $10^{30} \text{ cm}^{-3} \leq n_e \leq 10^{35} \text{ cm}^{-3}$.

At the density $n_e > 10^{35} \text{ cm}^{-3}$, the process of protonization of a nucleus becomes the main mechanism of nuclear transformations in the volume of a compressed Wigner–Seitz cell. In this interval, a decrease in the energy of the system at the expense of the increase in the binding energy of electrons with the nucleus becomes a more significant factor than the

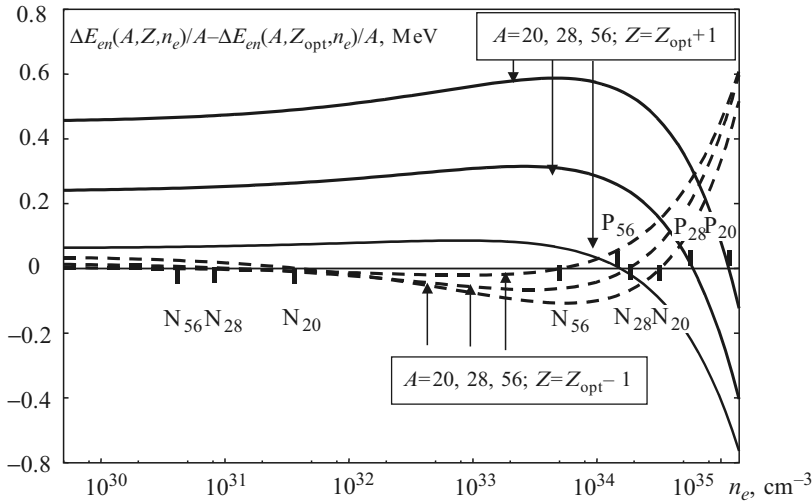


Fig. 11.14. Dependence of the total energy $\Delta E_{en}(A, Z, n_e)/A$ of a compressed quasiautom on the atomic number A and the density of the gas of degenerate electrons n_e on the line of stability Z_{opt} upon the deviation from it by $\Delta Z = -1$. Negative values of $\Delta E_{en}(A, Z_{opt}-1, n_e)/A - \Delta E_{en}(A, Z_{opt}, n_e)/A$ correspond to the neutronization region.

increase in the energy of the system at the expense of the increase in the Fermi energy of the compressed electron gas.

The comparison of the results presented in Figs. 11.11–11.14 yields that the process of neutronization runs at a relatively moderate density of the degenerate gas with decrease in both the density of electrons and the nucleus charge, whereas the process of protonization accompanied by the increase in the nucleus charge and the density of electrons runs at an extremely high density.

The above-derived results and the existence of the “neutronization corridor” raise the natural question: If the process of neutronization is a defining process of nuclear transformations in the region of electron densities $10^{30} \text{ cm}^{-3} \leq n_e \leq 10^{35} \text{ cm}^{-3}$, how can the density of the compressed electron gas exceeding the upper threshold of neutronization, $n_e \approx 10^{35} \text{ cm}^{-3}$, be reached? The obviousness of this question is related to that the region of preferred neutronization can be a distinctive filter, in the scope of which any increase in the density of electrons is accompanied by its automatic decrease (in the limit, up to the lower boundary of the region of neutronization). This circumstance can lead to such a situation where the region of protonization becomes unattainable!

One of the obvious answers to such a question can be related to that the energy gain of the process of neutronization does not mean that the neutronization is certainly realized. It is well known that the β -processes are connected with weak interactions, and the probability of such interactions is sufficiently low in all the cases. For this reason, the process of neutronization does not manage to occur for that small time, during which the density grows (e.g., at the expense of the mechanism of the collapse of an electron shell) in the scope of the whole “critical” interval of neutronization:

$$10^{30} \text{ cm}^{-3} \leq n_e \leq 10^{35} \text{ cm}^{-3}.$$

This argument becomes especially convincing if we take into account that the duration of the collapse is defined by the compression velocity of the electron shell which is close to the velocity of light at the last stage of the compression. In this case, the duration of the process of self-compression of the electron shell at the last, most important stage will not exceed 10^{-13} to 10^{-17} s in any case. It seems almost obvious that the probability of neutronization is very low during such small period. This leads to the very high probability of that the nucleus in the volume of a Wigner–Seitz cell at the stage of compression “passes across the critical region of neutronization” and falls into the region of protonization without change in its charge.

This conclusion is confirmed by a simple calculation.

Let us consider the quantitative aspect of the running of the reaction of neutronization and let us calculate its probability.

In the first order of perturbation theory, the probability of the process of neutronization (i.e., the reaction, in which a proton captures a certain electron, being in a state with energy E_e , is transformed into a neutron, and emits a neutrino with energy $E_\nu = E_e - E_n$) is described by the formula

$$P(E_e, E_p, E_n, E_\nu) = (2\pi/\hbar) |\langle E_e, E_p | H_{pn} | E_n, E_\nu \rangle|^2 \cdot \delta(E_e + E_p - E_n - E_\nu), \quad (11.80)$$

where the quantity $\langle E_e, E_p | H_{pn} | E_n, E_\nu \rangle$ is the matrix element of the process of neutronization, $p + e \rightarrow n + \nu$.

The sequential analysis of the weak interaction defining the transformation of a proton into a neutron on the basis of the Fermi theory shows that the value of this matrix element is connected with the constant of the weak interaction $g_s = 1.37 \times 10^{-49} \text{ erg cm}^3$ by the relation

$$|\langle E_e, E_p | H_{pn} | E_n, E_\nu \rangle|^2 = (g_s)^2 (|\int 1|^2 + R|\sigma|^2)/V^2. \quad (11.81)$$

Here, the quantity $(|\int 1|^2 + R|\sigma|^2)$ characterizes the normed dimensionless matrix elements for different types of the intranuclear transformations corresponding to the weak interaction with the participation of

a neutrino. This quantity will be different for different types of the transformations of nucleons, but it is of the order of several units for allowed β -transitions. For example, for the transformation of a free proton into a neutron (and a free neutron into a proton), the normed matrix element $(|\int 1|^2 + R|\sigma|^2) = 5$. The normalizing volume V corresponds to the region of existence (quantization) of the initial electron and the final neutrino.

Eq. 11.80 defines the probability of the one-channel transition by the indicated scheme: $E_e + E_p \rightarrow E_n + E_\nu$. However, the transformation $p + e \rightarrow n + \nu$ is multichannel (i.e., it can run with the participation of different electrons and, respectively, can lead to the creation of different neutrinos).

The total probability of the process of neutronization $p + e \rightarrow n + \nu$ can be determined by summation of Eq. 11.80 over all possible initial states of the system (over all electrons which are on different levels of energy E_e and E_p) and over all final states (over all states of a neutrino which satisfy the energy conservation law defining the argument of the delta-function) as

$$P_{p \rightarrow n} = \Sigma_e \Sigma_\nu P(E_e, E_p, E_n, E_\nu) \\ = \int_{E_n^{(\min)}}^{E_F + m_e c^2} \rho_e(E_e) dE_e \int_{E_n^{(\min)}}^{E_F + m_e c^2} \rho_\nu(E_\nu) dE_\nu P(E_e, E_p, E_n, E_\nu); \quad (11.82)$$

here,

$$E_e^{(\min)} = m_e c^2, \text{ if } E_n - E_p < m_e c^2, \quad (11.83)$$

$$E_e^{(\min)} = E_n - E_p, \text{ if } E_n - E_p > m_e c^2. \quad (11.84)$$

Now we introduce the densities of possible states of an electron $\rho_e(E_e)$ and a neutrino $\rho_\nu(E_\nu)$ in the scopes of the electron and neutrino Fermi-spheres as

$$\rho_e(E_e) = (d/dE_e) \left\{ 8\pi V \int_0^{E_e} p_e^2 dp_e / (2\pi\hbar)^3 \right\} \\ = (V/\pi^2 \hbar^3 c^3) E_e (E_e^2 - m_e^2 c^4)^{1/2}, \quad (11.85)$$

$$\rho_\nu(E_\nu) = (d/dE_\nu) \left\{ 8\pi V \int_0^{E_\nu} p_\nu^2 dp_\nu / (2\pi\hbar)^3 \right\} \\ = (V/\pi^2 \hbar^3 c^3) E_\nu^2. \quad (11.86)$$

After the substitution of Eqs. 11.81 and 11.86 in Eq. 11.82, we find the formula for the probability of the transformation (per unit time) of the given isotope into an isotope with the same mass number, but with a lesser charge of the nucleus (e.g., the transformation of a proton into a neutron):

$$P_{p \rightarrow n} = (2/\pi^3 \hbar^7 c^6)(g_s)^2 (|\int 1|^2 + R|\sigma|^2) \times E_e(E_e^2 - m_e^2 c^4)^{1/2} (E_e - E_n + E_p)^2 dE_e. \quad (11.87)$$

We will analyze this probability in two cases.

If the maximum of the electron energy slightly differs from $m_e c^2$ (i.e., when the Fermi energy $E_F < m_e c^2$), Eq. 11.87 yields

$$P_{p \rightarrow n} \approx (2/\pi^3 \hbar^7 c^6)(g_s)^2 (|\int 1|^2 + R|\sigma|^2)(m_e c^2)^3 E_F^2. \quad (11.88)$$

Substituting the formula for E_F , we get

$$P_{p \rightarrow n} \approx [g_s^2 n_e^{4/3} m_e / 2\pi^3 \hbar^3 (3\pi^2)^{4/3}] (|\int 1|^2 + R|\sigma|^2). \quad (11.89)$$

In the other limiting case where the Fermi level considerably exceeds $m_e c^2$, we get the following formula for the probability of the transformation of a proton into a neutron:

$$P_{p \rightarrow n} = (2g_s^2 / 5\pi^3 \hbar^7 c^6) (|\int 1|^2 + R|\sigma|^2) E_F^5. \quad (11.90)$$

In this case, a relation connecting the probability of the neutronization of a proton (in the both cases of a free proton and a proton in the composition of a nucleus) per unit time with a local density n_e of the electron gas in the volume of the nucleus is as follows:

$$P_{p \rightarrow n} = [2g_s^2 (3\pi^2)^{5/3} n_e^{5/3} / 5\pi^3 \hbar^2 c] (|\int 1|^2 + R|\sigma|^2). \quad (11.91)$$

If we take that such a neutronization is possible for each proton of a nucleus, then the final formula for the total probability of the neutronization of a nucleus containing Z protons looks

$$P_{Z(p \rightarrow n)} = Z P_{p \rightarrow n} = Z \left\{ 2g_s^2 (3\pi^2)^{5/3} n_e^{5/3} / 5\pi^3 \hbar^2 c \right\} (|\int 1|^2 + R|\sigma|^2). \quad (11.92)$$

The analysis of Eq. 11.92 shows that, other things being equal, the probability of the process of neutronization of nuclei increases sharply [proportionally to $P_{Z(p \rightarrow n)} \sim Z n_e^{5/3}$] with the increase of the product $Z n_e^{5/3}$. In result the increase of electronic density n_e is more important than the increase of a nuclear charge Z .

It was shown above (see Fig. 11.14) that for heavy nuclei with $A > 100$ the area of neutronization corresponds to lower level electron density $n_e \approx 10^{31} \text{ cm}^{-3}$ that for more light nuclei with $A < 60$ ($n_e \approx 10^{33}$ to $3 \times 10^{34} \text{ cm}^{-3}$). Using both this result and dependence Eq. 11.92, we can see that, for lighter nuclei, the probability of neutronization is much higher than that for heavy ones.

For example, in a plasma with the density of electrons $n_e = 10^{31} \text{ cm}^{-3}$, the probability of the single neutronization of a heavy nucleus with $Z = 92$ calculated on the basis of Eq. 11.92 is equal to $P_{92(p \rightarrow n)} \approx 10^4 \text{ s}^{-1}$. Finally, if the stay time in the “neutronization corridor” is $\Delta t = 10^{-13} \text{ s}$, we get that the probability of the single neutronization of this nuclei is very small:

$$W_{92(p \rightarrow n)} = P_{92(p \rightarrow n)} \Delta t \approx 10^{-6}. \quad (11.93)$$

For nuclei with lesser charges and lesser masses, this probability will be much higher. In particular, for nuclei of the “iron peak” with $Z \approx 26$ and $n_e = 10^{33} \text{ cm}^{-3}$, $P_{26(p \rightarrow n)} \approx 5 \times 10^8 \text{ s}^{-1}$, $W_{26(p \rightarrow n)} \approx 0.05$.

We note that, for such astrophysical objects as white dwarfs and a massive star on the early stage of the gravitational collapse, the velocity of compression will be considerably lesser (by virtue of the incomparably greater size of an object undergoing the compression). In particular, if we take that $\Delta t \gg 10^{-7} \dots 10^{-6} \text{ s}$ (see Ref. 44), it is easy to verify that the total probability of the neutronization will be close to 1.

This circumstance can be that main reason which does not “allow” the nuclei in astrophysical objects “to leave” the scope of the “neutronization corridor” for the region with a higher electron density and to realize the process of protonization.

In the system under consideration (in the case of a small-size target), the process of compression can run so rapidly that the probability of the neutronization turns out to be low. In this case, the nuclei rapidly traverse the “neutronization corridor” and come to the region with the electron density corresponding to the process of protonization.

The character of the running of nuclear processes at higher electron densities is related to many acts of the creation of heavier particles and requires the additional study.

11.2. Evolution of Self-Controlled Electron-Nucleus Collapse in Condensed Targets and a Model of Scanning Nucleosynthesis

11.2.1. Stability of Matter and the Problem of Collapse under Laboratory Conditions

The general problem of stability of the matter together with the problem of collapse of the substance of an electron-nucleus system, problem of realization of global electron-nucleus transformations, and problem of controlled energy generation which follow from the first one belong to the most fundamental trends of modern science. Earlier, the problem of stability of a substance was considered only in connection with such global phenomena of astrophysics as a gravitational collapse, the formation of neutron stars, and bursts of supernovas (see Ref. 44). The possibility to realize the substance collapse under laboratory conditions was not analyzed at all due to the absence of available experimental methods. In recent years, this problem was most closely approached by the physics of inertial thermonuclear synthesis, in the framework of which the increase in the target density by 100 to 200 times is considered to be a necessary condition of the energy-gained synthesis. It is clear, however, that the multiple increase in the target density at the expense of the action of a very powerful external driver has no direct relation, by a high account, to the process of collapse of a substance. It is one of the necessary conditions, but it is far from being a sufficient one. The substance collapse can be realized, in our opinion, only at the expense of self-similar processes related to the action of internal mechanisms weakly depending on the initial conditions and the initial action of an external driver. The proper collapse should be a self-organizing process. In fact, a collapsing system “forgets” its prehistory very fast and is developing by the own “inner” laws. The reason for this is related to that the action of internal mechanisms of self-compression of the matter is incommensurably more essential than that of an initial perturbation which is, in fact, a distinctive trigger.

Recently, it has been convincingly demonstrated that the state of collapse can be also realized under conditions of a terrestrial laboratory under the energy contribution to be at most 1 kJ. This new trend was theoretically substantiated and confirmed by numerous experiments that were carried out at the Electrodynamics Laboratory “Proton-21” in Kiev (Refs. 53, 52, 55, 56).

Very important is the circumstance that, even at a relatively small energy contribution, the state of collapse involves a macroscopic amount of the matter (at least 10^{20} nucleons), rather than two individual nuclei.

The latter situation was considered by the projects on the use of high-energy accelerators for the creation of a quark-gluon plasma in the head-on collisions of accelerated nuclei.

The detailed studies (see Ref. 56) have shown that the phenomenon of global self-compression of the matter can be related to the Coulomb collapse of an electron-nucleus system. The conditions for such a collapse can be satisfied in the case where a macroscopic amount of a substance is preliminarily compressed to the critical density which corresponds to a degenerate relativistic electron gas and depends on the charge of nuclei. The calculations (see Ref. 56) yield that the minimum specific energy introduced into the compressed plasma to ensure the conditions of a collapse should be equal to 0.6 to 2 MeV/electron. In this case, the minimum energy corresponds to the heaviest nuclei (nuclei with maximum charge Z). This condition can be met naturally in astrophysical phenomena during the preceding gravitational collapse of a heavy star, but it can hardly be satisfied under conditions of a usual laboratory by using the traditional methods of compression of targets.

Below we consider several different methods of compression of a substance to the critical density. One of them is connected with the use of a large pressure,

$$P_J = H_\varphi^2/8\pi, \quad (11.94)$$

induced by the azimuthal magnetic field $H_\varphi = 2J/Rc$ that appears in a micropinch system under the pulse compression of the channel of a current J to that with the radius R determined from the condition of the equilibrium between the kinetic pressure of the electron gas and P_J . Our estimations showed that, with regard to inertial forces, a very little amount of a substance can reach the state of degenerate relativistic electron gas with the density sufficient for the start of a collapse during the pulse compression of a micropinch in the condensed medium. However, the analysis shows that the collapse of a system on the basis of a magnetic micropinch has a very low efficiency and gives no foundation for the explanation of the results of numerous successful experiments carried out since 1999 at the Electrodynamics Laboratory "Proton-21" (see Refs. 52, 55, 56):

1. Due to the small energy of the initial electron beam (at most 1 kJ), the amount of a compressed substance in the zone of action of a micropinch is also very small and does not exceed 10^{14} to 10^{15} electrons and nucleons. At the same time, we observed a nuclear transformation of the substance involving more than $N_1 \approx 10^{20}$ nucleons in each of the performed experiments.

2. The channel of the current micropinch has the form of a very thin (at most $1\ \mu\text{A}$ in diameter) long thread, and the zone of action of the collapse in experiments corresponds always to a macroscopic sphere with radius exceeding $50\ \mu\text{m}$.
3. A significant part of the products of nuclear transformations is a cold condensed substance (hardened drops or monolithic sufficiently thick layers) in the region adjacent to the collapse zone, rather than various atoms and molecules randomly distributed over the volume or surface.

We developed the other (stage-by-stage or scanning) conception of the formation of a collapse which asserts that the initial state of collapse is formed in a thin spherical layer containing ΔN electrons neutralized by ions (by ΔN protons of nuclei). This layer moves with acceleration to the target center by rapidly increasing the own density and energy (see Ref. 56). Upon the movement of this layer in the unperturbed solid target, there occur the compression of the ion (nuclear) component of a substance, neutralization of the charge of electrons, and formation of the electron-nucleus collapse on the leading edge of the layer. On the trailing edge of the spherical layer moving to the target center, there occur the breaking of the state of electron-nucleus collapse, partial restoration of the target structure, and related fast adiabatic cooling of the products of nuclear transformations, which explains the observed peculiar localization of synthesized elements and isotopes.

It follows from this conception that the process of nuclear transformations involves, by turns, the whole volume of the substance ($N_1 \gg \Delta N$ nucleons), through which the spherical plasma layer undergoing compression to the state of collapse, i.e., a self-controlled “wave-shell,” moves. The action of the collapse in the form of an extremely compressed propagating thin layer is terminated at the target center by the formation of a spherical collapse, its inertial holding, a number of fundamental nuclear transformations, and the subsequent decay. This conception allows us to combine two, as if mutually contradictory, facts: the small energy contributed, in the initial state, to this high-density wave in the target and the great volume of a substance undergoing a global transformation in the process of action of the wave.

Below, we will first show that solving the problem of realization of a nonstationary self-compression of a part of the target in the form of a spherical layer terminating in the state of collapse can be related to peculiarities of the interaction of the bounded system of quasifree electrons with the ion matrix of the target. In this case, the formal start of the process of self-compression does not require a preliminary supercritical action.

In Sec. 11.2.2, we substantiate the potentialities and efficiency of such a process by basing on the detailed analysis of peculiarities of the interaction

of the bounded system of a degenerate electron gas with the multiply ionized target. In Sec. 11.2.3, we study the general regularities of the evolution of a degenerate Fermi-gas of electrons in condensed targets in the presence of a drift motion of electrons. Peculiarities of the formation and motion of a scanning spherical layer of a degenerate electron gas in a condensed target are considered in Sec. 11.2.4 Section 11.2.5 is devoted to the study of peculiarities of the motion of the ion (nuclear) component of a target in the volume of a scanning spherical layer. In this section, a particular attention is paid to the analysis of the mutual influence of the electron and ion components of a target on the motion of a scanning high-density plasma layer. In the same section, we analyze comprehensively the premises of the formation and the peculiarities of the propagation of a shock wave synchronized with the movement of the scanning plasma layer.

This chapter ends with both the analysis of general regularities of the “scanning” synthesis and the detailed analysis of the spatial distribution of synthesized isotopes in accumulating screens remote from the collapse zone.

11.2.2. Interaction of the Bounded System of a Degenerate Electron Gas with a Multiply Ionized Target

Consider the general regularities of the interaction of a degenerate electron gas with the multiply ionized target. We assume that a part of atoms has been ionized in a thin spherical layer near the surface of the condensed target with spherical symmetry under the action of an external perturbation. Such an ionization can be caused by a shock wave, being a consequence of the pulse external heating or a correlated intense short-time impact on the surface. The other version of the formation of a thin ionized layer can be related to the introduction of a bounded part of the electron beam, which forms a propagating electron layer, into the target from the inside along a normal to the surface. Other methods to generate the initial perturbation in the near-surface layer of the target can be considered as well.

In the case under study, we assume that, during a certain time interval (not exceeding the time interval of the inverse recombination of electrons and ions), the target contains a free electron gas with density n_e and the totality of partially ionized atoms (ions) neutralizing the charge of electrons. If the temperature of this gas is less than the Fermi energy, the metallic coupling of these electrons and ions will be realized in the scope of such a system.

For example, if a target is a dielectric, it becomes a metal. If it is already a metal, its metallic properties will be considerably enhanced.

The total energy of the electron–nucleus interaction in such a system will be defined by the interaction of the degenerate electron Fermi-gas with ions.

The various components of this energy were analyzed earlier (see Ref. 56) and considered above in Chapter 11.1. (see Eqs. 11.22–11.34).

In addition, because this electron gas has a relatively small density (as compared to that of a superrelativistic degenerate gas which is formed at the last stages of the analyzed process of self-compression), the beginning of the process of compression will considerably depend on correlation effects leading to a partial screening of the field of ions. This last effect can be taken into account by using, for example, the interpolational formula of Nozières and Pines (see Ref. 61)

$$U_{\text{corr}} = -(e^4 m / \hbar^2)(0.115 - 0.031 \ln r_s) \quad (11.95)$$

defining the correlation energy in the region $2 < r_s < 5$ corresponding to all real metals.

Here,

$$r_s = R/a_B \equiv (me^2/\hbar^2)(4\pi/3)^{1/3}Z^{1/3}/n_e^{1/3}, \quad (11.96)$$

$a_B := \hbar^2/me^2$ is the radius of the Bohr orbit, and $R = (3/4\pi)^{1/3}Z^{1/3}/n_e^{1/3}$ is the radius of a Wigner–Seitz cell determined from the condition $(4\pi/3)R^3n_e = Z$ for an ion with charge Z .

This formula describes satisfactorily the correlation effects for the electron gas density

$$\begin{aligned} (3/4\pi)Z/(a_B r_{s(\text{max})})^3 &= 2 \times 10^{22} Z \text{ cm}^{-3} \\ &< n_e < \\ (3/4\pi)Z/(a_B r_{s(\text{min})})^3 &= 2 \times 10^{23} Z \text{ cm}^{-3}. \end{aligned} \quad (11.97)$$

At a higher density of the degenerate electron gas (all the more, in the case of a degenerate relativistic electron gas), the influence of correlation effects becomes insignificant. Finally, using Eqs. 11.22–11.34 and Eq. 11.95, we find the total energy of the electron subsystem of a Wigner–Seitz cell as

$$\begin{aligned} U_{e\Sigma} &= U_{eQL} + U_{\text{corr}} + U_{eQNL} + U_{e,exch} + U_{eF} \\ &= [W_{Q(en)} + U_{\text{corr}} + W_{Q(ee)} + Zm_e c^2] K_F \\ &\quad + U_{e,exch} - Zm_e c^2 + U_{eQNL} \\ &= \left\{ Zm_e c^2 - (9/10)(4\pi/3)^{1/3}Z^{5/3}e^2n_e^{1/3} \right. \\ &\quad \left. + (e^4 m / \hbar^2)(-0.115 + 0.031 \ln[(e^2 m / \hbar^2)(4\pi/3)^{1/3}Z^{1/3}/n_e^{1/3}]) \right\} K_F \\ &\quad - [3(3\pi^2)^{1/3}/4\pi]Ze^2n_e^{1/3} \\ &\quad - Zm_e c^2 + [8\pi Z^{7/3}e^4(3/4\pi)^{1/3}/7m_e c^2]n_e^{2/3}. \end{aligned} \quad (11.98)$$

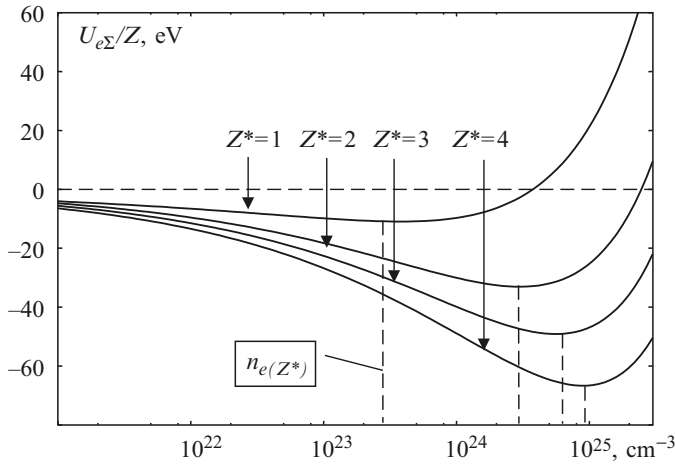


Fig. 11.15. Interaction energy between a low-charge ion and degenerate electrons in a Wigner–Seitz cell.

In Fig. 11.15, we present this energy as a function of the ion charge (for ions with charges $Z^* = 1, 2, 3$, and 4) and of the density of the electron gas surrounding ions.

As seen, the increase in the multiplicity of ionization of an atom, Z^* , induces two main effects:

1. A systematic displacement of the position of a minimum of the total energy $U_{e\Sigma}$ of the electron subsystem of a Wigner–Seitz cell, which defines the stable state of this system, to the side of a greater density of electrons (and, as a result, of a greater density of ions) in the spatial region including the system of Wigner–Seitz cells)
2. A sharp and systematic deepening of this minimum (the depth of the potential well) defining the position of a stable state of the system of electrons and ions of the corresponding multiplicity

The derived results yield that the position, n_{eZ^*} , and the depth of a minimum of the total energy of the electron subsystem of a Wigner–Seitz cell depend on the ion multiplicity Z^* . In particular, at $Z^* = 1$, we have $n_{e(Z^*=1)} \approx 10^{23} \text{ cm}^{-3}$, which corresponds to the typical concentration of electrons under the metallic type of coupling in any condensed target. The equilibrium concentration of electrons n_{eZ^*} rapidly increases with Z^* : at $Z^* = 2, 3$, and 4 , we have, respectively, $n_{e(Z^*=2)} \approx 3 \times 10^{23} \text{ cm}^{-3}$, $n_{e(Z^*=3)} \approx 6 \times 10^{24} \text{ cm}^{-3}$, and $n_{e(Z^*=4)} \approx 10^{25} \text{ cm}^{-3}$.

The same regularity (a sharp displacement of the position of the energy minimum and an increase in its depth) will occur with increase in

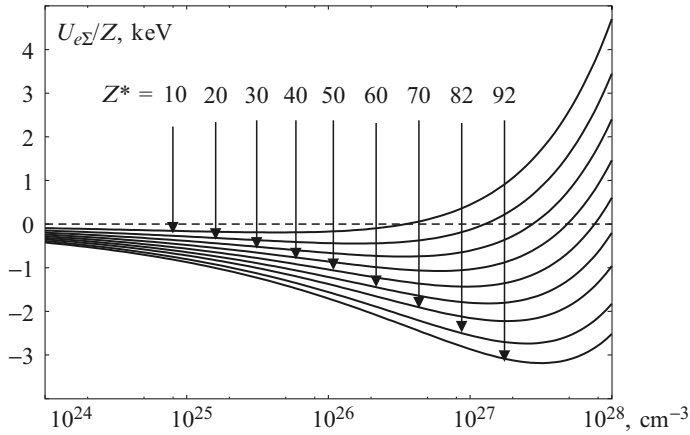


Fig. 11.16. Interaction energy between ions with a high charge and degenerate electrons in a Wigner-Seitz cell.

the degree of ionization of heavy atoms, i.e., upon the further increase in Z^* (Fig. 11.16).

For example, for iron with $Z^* = 26$, we have $n_{e(Z^*=26)} \approx 3 \times 10^{26} \text{ cm}^{-3}$, and $n_{e(Z^*=92)} \approx 4 \times 10^{27} \text{ cm}^{-3}$ for the heaviest stable nucleus (the uranium nucleus with $Z^* = 92$).

The mentioned effects yield the following very important circumstance.

It follows from the general principles of thermodynamics that any nonequilibrium system relaxes to a stable state defined by the maximum of the binding energy (or the minimum of the total energy). Let us assume that we create a system with a specific degree of ionization Z^* such that its electron density n_e is less than the optimum equilibrium value $n_{e(Z^*)}$ at the moment of the creation. Then we can observe the effect of the fast self-compression of a region in the target under consideration up to this optimum density. This process must run very rapidly because it is a consequence of the action of the very great forces induced by the violation of the balance of Coulomb attraction and Fermi repulsion in every Wigner-Seitz cell.

For the sake of justice, we note that there exist other processes which can compete with the self-compression of a target upon the violation of its charge state.

In particular, if the linear size of a region with primary ionization is very large, the considered process of global “metallization” will last longer than the reverse recombination of electrons. In this case, the recombination can be realized faster than the Fermi collective state of the target will be

formed. The notion of “large linear size” means that the time, for which a Fermi electron passes the entire region, will be greater than the recombination time. In this case, the recombination occurs faster than the formation of a Fermi-condensate. In such a system, the global self-compression is impossible.

But if the linear size of the region of primary ionization (e.g., the thickness of the spherical layer of a primary perturbation) is small, the Fermi collective interaction will be formed prior to the termination of the recombination. In this case, naturally, the recombination of electrons into the old “atomic state” is impossible, and the process of self-compression becomes the defining factor of the evolution.

It is easy to carry out the corresponding quantitative estimations.

In such a system, the main mechanism of recombination is the photorecombination, whose efficiency and duration τ_φ are defined by the probability of a spontaneous radiative transition of an electron, which was formed due to the ionization, into a bound atomic state. For the majority of elements, the ionization potentials φ_i of the first and second electrons correspond to the near UV range and do not exceed $\varphi_{i(1)} \approx 5$ to 10 eV and $\varphi_{i(2)} \approx 20$ to 30 eV. The photorecombination duration $\tau_\varphi = 3\hbar c^3 / 4\omega^3 |d_{in}|^2$ for allowed dipole transitions with $\omega \approx \varphi_i / \hbar$, $|d_{in}| \approx er_B \approx \hbar^2 / m_e e$ is equal to $\tau_\varphi \approx 0.1$ to 1 ns. It follows from these estimations that, for the typical velocity of Fermi electrons in the conduction zone of a metal, $v_F \approx 10^8$ cm/s, the maximum size of the region of primary ionization (in the model under study, it is the thickness of a spherical plasma layer ΔR near the target surface), for which the effect of metallization of the matrix is possible, does not exceed $\tau_\varphi v_F \approx 100$ to 1000 μm .

11.2.3. Evolution of a Degenerate Fermi-Gas of Electrons in Condensed Targets in the Presence of a Drift Motion of Electrons

The above-considered peculiarities of a deformation of the target upon a change of the charge state of its electron-nucleus system allow us to raise a quite natural question: Is it possible to realize the process of chain ionization and self-compression of a target? Such a process corresponds to the situation where the self-compression of a medium, in which the fast ionization (with the multiplicity of ionization Z^*) has happened, from the initial state with electron density n_e to a stable state with higher electron density $n_{e(Z^*)} > n_e$ corresponding to a specific value of Z^* , is associated with such a deformation of the medium which leads to the increase in the ionization multiplicity up to $(Z^* + \Delta Z^*)$. The last corresponds, in turn, to a higher optimum density $n_{e(Z^* + \Delta Z^*)}$.

It is obvious that if such an effect is possible in a certain condensed target, the process of ionization and self-compression of this target acquires the avalanche-like character. To this end, it is sufficient to create a relatively small initial ionization (with small Z^*) in the target with $Z^* = 0$. Then the target substance begins to compress itself up to the equilibrium density $n_{e(Z^*)}$ corresponding to Z^* and, in this case, ionizes itself to $Z^* + 1$. The new value of ionization multiplicity $Z^* + 1$ leads to the effect of a further compression to a new equilibrium density $n_{e(Z^*+1)}$. This leads, in turn, to the increase in the ionization multiplicity up to $Z^* + 2$, and so on.

As a result of such an avalanche-like multiple mutual stimulation of the processes of ionization of the target atoms and deformation (self-compression) of the target, its charge state attains rapidly the characteristics corresponding to the maximum (full) ionization with $Z^* = Z$. In this case, the completely ionized target will be a plasma compressed to the maximum density $n_{e(Z)}$.

The simple estimations showed that such a process is impossible under the "usual" conditions typical of a condensed target (e.g., a solid body). To make such an assertion, it is sufficient to determine the density of electrons which corresponds to the breaking of the next atomic shell and the collectivization of its electrons. This threshold density can be determined from the condition, according to which the external shell of any ion with charge $Z^* \leq Z$ is completely broken, when the ion density reaches the threshold value (see Ref. 62) $n_{icr} \approx 0.5Z^*(m_e e^2 / \hbar^2)^3 \approx 0.5 \times 10^{25} Z^* \text{ cm}^{-3}$.

This value of n_i corresponds to the threshold electron density $n_{ecr} \approx Z^* n_{icr} \approx 0.5 \times 10^{25} (Z^*)^2 \text{ cm}^{-3}$. For the deepest levels of multiply ionized ions with $Z^* \approx Z > 70$ to 80, the threshold density of the breaking of a shell is $n_{ecr} \approx 10Z^* n_{icr} \approx 0.5 \times 10^{26} (Z^*)^2 \text{ cm}^{-3}$.

This condition yields that the thresholds of the formation of ions with $Z^* = 2, 3$, and 4 are equal to, respectively,

$$\begin{aligned} n_{ecr(Z^*=2)} &\approx 2 \times 10^{25} \text{ cm}^{-3}, \\ n_{ecr(Z^*=3)} &\approx 5 \times 10^{25} \text{ cm}^{-3}, \text{ and} \\ n_{ecr(Z^*=4)} &\approx 8 \times 10^{25} \text{ cm}^{-3}. \end{aligned}$$

An analogous situation will hold for heavier atoms.

In the case of the full ionization of Fe with $Z^* = Z = 26$, we get $n_{ecr(Z^*=26)} \approx 3 \times 10^{27} \text{ cm}^{-3}$. For the full ionization of U ($Z^* = Z = 92$), $n_{ecr(Z^*=92)} \approx 4 \times 10^{29} \text{ cm}^{-3}$, which is larger by 10 to 100 times than the equilibrium electron density n_{eZ^*} for the same ion.

It is seen from these estimations that the concentration of electrons $n_{e(Z^*)}$ in an ionized substance, which corresponds to a stable state of an ion with charge Z^* , is significantly less than the critical electron concentration

$n_{ecr}(Z^*)$, at which an ion with a higher charge $Z^* + 1$ is formed due to the self-compression of a target (e.g., of the metallic matrix). This result would be expected, because any substance would be extremely unstable in the opposite case: the initial action of an arbitrary small initial deformation or the initial ionization would be sufficient in order that the process of self-compression transform this substance in a completely ionized superdense electron-nucleus degenerate plasma. Nobody has ever observed stationary media which possess both the ability of an avalanche-like ionization and such anomalous properties of the instability. Any condensed body (except for bodies with mass exceeding the threshold of gravitational collapse) is absolutely stable as for the process of unbounded compression and admits a bounded compression only up to the state of internal equilibrium. As known, the latter is realized, for example, in any metal.

The reason for such a stability of metals is well known and is conditioned by that their compression occurs only up to such a limiting electron density $n_{e(Z^*)}$, at which the attraction of electrons to nuclei in the volume of every Wigner–Seitz cell is balanced by the kinetic (Fermi) pressure of the degenerate electron gas.

The analysis indicates that, despite the seeming obviousness of such a conclusion on the absolute stability of substance, it is not true in all the cases, because it is derived under definite conditions. In particular, the above-presented condition of equilibrium characterizes the state of a Fermi-gas with zero mean momentum of electrons, i.e., in the absence of the ordered drift motion (though the square mean momentum of electrons of a Fermi-gas turns out to be very high and is defined by the Fermi energy). In the static case, the condition of equilibrium corresponds to only one definite stable state with $n_{e(Z^*)}$ uniquely defined by the ion charge Z^* . It is easy to verify that, in dynamic systems (in particular, in the presence of a nonzero mean momentum of the drift motion $\vec{p}_0$), the condition of equilibrium can be controlled and can be displaced to the side of large values of $n_{e(Z^*)}$. Under certain conditions, such an effect of the parametric coupling between the drift momentum $\vec{p}_0$ and $n_{e(Z^*)}$ can ensure the fulfillment of the above-mentioned criterion of instability of a metal relative to the process of self-compression and, for this reason, can induce a fast transformation of any substance in a completely ionized and strongly compressed electron-nucleus plasma!

Consider the following model. Let us have a degenerate electron gas with density n_e . If nuclei neutralizing this gas are in a maximally stable (ordered) state and form a cubic lattice, then the quasicontinuous distributions of momenta and energy are set in every Wigner–Seitz cell by the Fermi distribution with a definite limiting momentum p_F .

What is happen if the motion of all electrons will be characterized by the simultaneously realized values of both the quascontinuously distributed momenta $\vec{p}$, whose directions and values correspond to a degenerate Fermi-gas, and a drift momentum $\vec{p} = \gamma m \vec{v}_0$?

In the associated coordinate system (moving with a velocity $\vec{v}_0$), electrons occupy all energy levels from $p_{\min} = 0$ to $p_{\max} = p_F = (3\pi^2 \hbar^3 n_e)^{1/3}$. In the laboratory system in the case of a drift in one direction, the Fermi-sphere is shifted by the value of the drift momentum $\vec{p}_0$ (Fig. 11.17a).

A similar phenomenon (but on an incommensurably smaller scale) occurs with conduction electrons in the presence of an electric field applied to the conductor.

In what follows, we will consider the real process of symmetric radial compression of a spherical layer. In this case, the degenerate electrons in the volume of the spherical layer have, besides the randomly oriented Fermi components of a momentum $\vec{p}_0 = \gamma m v_0 \vec{e}_r$, the ordered radial component $\vec{p}$ which is directed to the center of the system and is the same in value for all angles θ and φ (Fig. 11.17b).

Consider the case where the drift momentum exceeds the initial radius of the Fermi-sphere, i.e., $p_0 > (3\pi^2 \hbar^3 n_e)^{1/3}$. In the laboratory system, the momenta of the degenerate electrons in the scope of a Fermi-layer of $\Delta p_F(\theta, \varphi, p_0)$ in thickness are distributed in the interval from

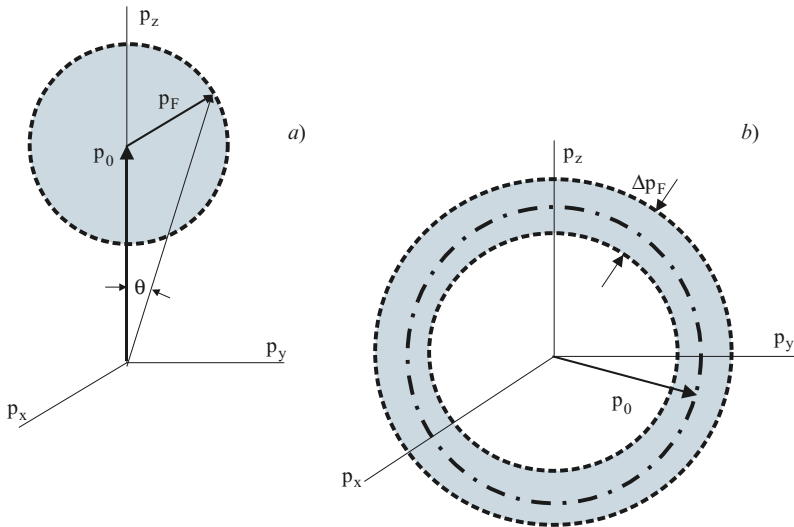


Fig. 11.17. Shift of the Fermi-sphere (a) and the formation of a spherical Fermi-layer (b) in the presence of unidirectional (a) and spherically symmetric (b) drift momenta p_0 of the gas of degenerate electrons.

$p_0(\theta, \varphi) - \Delta p_F(\theta, \varphi, p_0)/2$ to $p_0(\theta, \varphi) + \Delta p_F(\theta, \varphi, p_0)/2$. In the general case, the total number of occupied states in the volume V is defined by the relation

$$N_F = 2 \int d^3x \int d^3p / (2\pi\hbar)^3$$

$$= 2V \int_0^{2\pi} d\varphi \int_0^\pi \sin\vartheta d\vartheta \int_{p_0(\theta, \varphi) - \Delta p_F(\theta, \varphi, p_0)/2}^{p_0(\theta, \varphi) + \Delta p_F(\theta, \varphi, p_0)/2} p^2 dp / (2\pi\hbar)^3. \quad (11.99)$$

Consider the process of symmetric radial compression of a spherical layer with volume $V = N_0 V_0$, where N_0 nuclei and $N_F = ZN_0$ degenerate electrons are present.

Here, $V_0 = V/N_0$ is the volume of a Wigner–Seitz cell.

In this case, the electrons in the volume of the spherical layer have, besides the randomly oriented Fermi components of a momentum, the ordered radial component $\vec{p}_0 = \gamma m v_0 \vec{e}_r$ which is directed to the center of the system and is the same in value for all spatial angles θ and φ .

For such a system, we have $p_0(\theta, \varphi) \equiv p_0$, $\Delta p_F(\theta, \varphi, p_0) \equiv \Delta p_F(p_0)$. The considered physical model corresponds to the plasma layer that has a high density of degenerate electrons and moves to the center of the target (Fig. 11.18).

At any instantaneous position of this layer, there occurs the process of displacement of nuclei of the target in the layer in order to neutralize the electric charge of the gas of degenerate electrons. In this case, a deformation of the nuclear component in the form of a shock wave will be most probable. This question will be analyzed below in detail. It is obvious that different values of the layer radius R will correspond to different nuclei in the volume of the plasma layer.

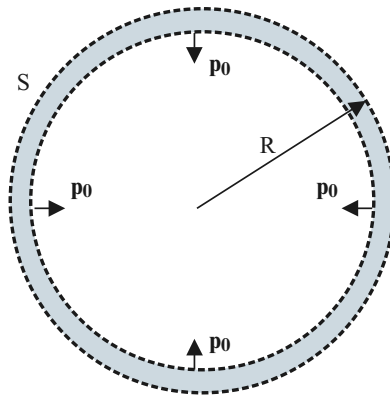


Fig. 11.18. Spherical plasma layer with a high electron-nucleus density moving to the center of a condensed target.

The total number of states for $N_F = ZN_0$ electrons in the spherical layer compressing itself is

$$2V \int_0^{2\pi} d\varphi \int_0^{\pi} \sin \vartheta d\vartheta \int_{p_0 - \Delta p_F(p_0)/2}^{p_0 + \Delta p_F(p_0)/2} p^2 dp / (2\pi\hbar)^3$$

$$= N_0 V_0 [p_0^2 \Delta p_F(p_0) + (\Delta p_F(p_0))^3 / 12] / \pi^2 \hbar^3. \quad (11.100)$$

This formula yields both the coupling equation for the concentration of degenerate electrons $n_e = Z/V_0$ and the limiting momentum

$$(\Delta p_F(p_0))^3 + 12p_0^2 \Delta p_F(p_0) = 12\pi^2 \hbar^3 n_e \quad (11.101)$$

and the formula for the limiting momentum

$$p_F(p_0) = p_0 + \Delta p_F(p_0)/2, \quad (11.102)$$

$$\Delta p_F(p_0) = \left\{ 6\pi^2 \hbar^3 n_e + [(6\pi^2 \hbar^3 n_e)^2 + 64p_0^6]^{1/2} \right\}^{1/3}$$

$$+ \left\{ (6\pi^2 \hbar^3 n_e) - [(6\pi^2 \hbar^3 n_e)^2 + 64p_0^6]^{1/2} \right\}^{1/3}. \quad (11.103)$$

Consider two extreme cases (relative to small and large drift velocities):

(a) *The case where the energy of the drift motion of the electron gas slightly exceeds its Fermi energy.*

In the case where the drift momentum p_0 is approximately equal to $\Delta p_F/2$ [i.e., if $p_0 \geq (3\pi^2 \hbar^3 n_e)^{1/3}$], relation Eq. 11.99 yields

$$p_F = (3\pi^2)^{1/3} \hbar n_e^{1/3} + p_0^3 / 3(3\pi^2)^{2/3} \hbar^2 n_e^{2/3}. \quad (11.104)$$

This value of p_F corresponds to the limiting (Fermi) energy

$$E_F = p_F^2 / 2m = (3\pi^2)^{2/3} \hbar^2 n_e^{2/3} (1 + 2p_0^3 / 9\pi^2 \hbar^3 n_e) / 2m, \quad (11.105)$$

for a nonrelativistic degenerate electron gas, and

$$E_F = p_F c = (3\pi^2)^{1/3} \hbar c n_e^{1/3} (1 + p_0^3 / 9\pi^2 \hbar^3 n_e) \quad (11.106)$$

for a superrelativistic gas.

(b) *The case where the energy of the drift motion of the electron gas is considerably more than the Fermi energy of this gas.*

In this limiting case [at a large drift velocity and, respectively, a large drift momentum $p_0 \gg (3\pi^2\hbar^3 n_e)^{1/3}$], Eq. 11.103 yields

$$p_F(p_0) = p_0[1 + \pi^2\hbar^3 n_e/2p_0^3]. \quad (11.107)$$

This value of p_F corresponds to the Fermi energy

$$E_F = p_F c = p_0 c[1 + \pi^2\hbar^3 n_e/p_0^3] \quad (11.108)$$

for a superrelativistic gas.

The total kinetic energy of all electrons in the volume of one cell is

$$\begin{aligned} U_{eF} &= \int_{p_0 - \Delta p_F(p_0)/2}^{p_0 + \Delta p_F(p_0)/2} (V_0/\pi^2\hbar^3) [(p^2 c^2 + m_e^2 c^4)^{1/2} - m_e c^2] p^2 dp \\ &= Z m_e c^2 [K_F(p_1, p_2) - 1], \end{aligned} \quad (11.109)$$

where

$$\begin{aligned} K_F(p_1, p_2) &= \int_{p_0 - \Delta p_F(p_0)/2}^{p_0 + \Delta p_F(p_0)/2} (V_0/Z m_e c^2 \pi^2 \hbar^3) (p^2 c^2 + m_e^2 c^4)^{1/2} p^2 dp \\ &= (m_e^3 c^3 / 8\pi^2 \hbar^3 n_e) \{ \{ (p_2/m_e c) [2(p_2/m_e c)^2 + 1] \\ &\quad \times [(p_2/m_e c)^2 + 1]^{1/2} - \text{Arsh}(p_2/m_e c) \} \\ &\quad - \{ (p_1/m_e c) [2(p_1/m_e c)^2 + 1] \\ &\quad \times [(p_1/m_e c)^2 + 1]^{1/2} - \text{Arsh}(p_1/m_e c) \} \}, \end{aligned} \quad (11.110)$$

$$p_{2,1} = p_0 \pm \Delta p_F(p_0)/2. \quad (11.111)$$

The presence of the ordered radial (drift) momentum $\vec{p}$ changes also the exchange energy of the degenerate electron gas defined, in this case, by the formula

$$\begin{aligned} U_{e,exch} &= -(4\pi Z \hbar^2 e^2 / n_e) \int_{p_0 - \Delta p_F(p_0)/2}^{p_0 + \Delta p_F(p_0)/2} d^3 p_1 / (2\pi\hbar)^3 \int_{p_0 - \Delta p_F(p_0)/2}^{p_0 + \Delta p_F(p_0)/2} d^3 p_2 / (2\pi\hbar)^3 |\vec{p}_1 - \vec{p}_2|^2 \\ &= -(Z e^2 / 4\pi^3 \hbar^4 n_e) \{ [(p_0 + \Delta p_F(p_0)/2)^4 - (p_0 - \Delta p_F(p_0)/2)^4] \\ &\quad - (4/3)(p_0 - \Delta p_F(p_0)/2) \\ &\quad \times [(p_0 + \Delta p_F(p_0)/2)^3 - (p_0 - \Delta p_F(p_0)/2)^3] \}. \end{aligned} \quad (11.112)$$

It follows from these relations that, with increase in the drift momentum p_0 , the exchange energy decreases, and its influence becomes insignificant upon a very fast motion.

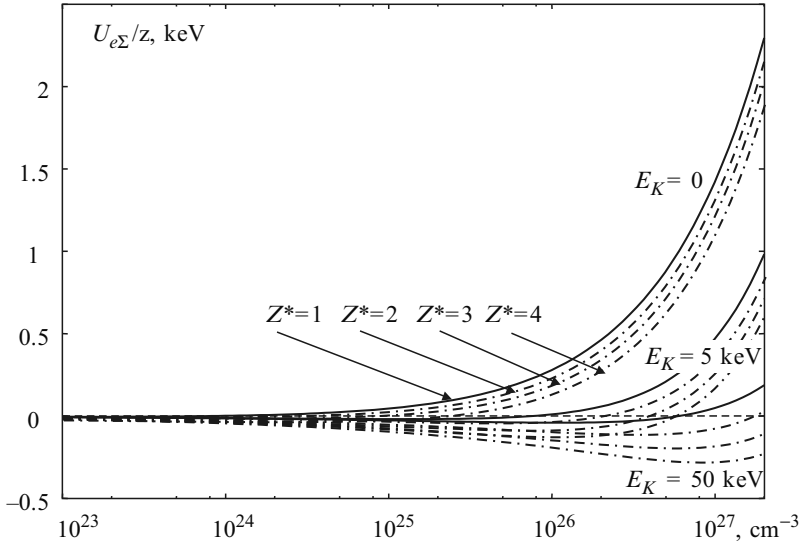


Fig. 11.19. Numerical analysis of the change in the total energy of a degenerate electron gas in one Wigner–Seitz cell for the low charges of ions Z^* as a function of the kinetic energy of drift motion of electrons, $E_K = (p_0^2 c^2 + m_e^2 c^4)^{1/2} - m_e c^2$.

In Fig. 11.19, we present the result of the numerical analysis of the total energy of a degenerate electron gas for the same charges of ions Z^* as a function of the kinetic energy of the drift motion of electrons $E_K = (p_0^2 c^2 + m_e^2 c^4)^{1/2} - m_e c^2$. These results are derived on the basis of the general formula Eq. 11.98 modified with regard to the dependence of the energy U_{eF} given by Eq. 11.109 on the drift momentum p_0 .

It is seen from this figure that, already at a relatively small energy of the drift motion $E_K = 5 \text{ keV}$, the energy minimum (the stable state of the electron-ion system) shifts from the electron densities $n_{e(Z^*=2)} \approx 5 \times 10^{23} \text{ cm}^{-3}$, $n_{e(Z^*=3)} \approx 10^{24} \text{ cm}^{-3}$, and $n_{e(Z^*=4)} \approx 4 \times 10^{24} \text{ cm}^{-3}$ characteristic of a static degenerate gas of electrons to $n_{e(Z^*=2)} \approx 4 \times 10^{25} \text{ cm}^{-3}$, $n_{e(Z^*=3)} \approx 8 \times 10^{25} \text{ cm}^{-3}$, and $n_{e(Z^*=4)} \approx 1.3 \times 10^{26} \text{ cm}^{-3}$. As seen, the presence of the drift motion with a relatively small energy leads to the increase in the equilibrium electron density $n_{e(Z^*)}$, $Z^* = 2, 3, 4, \dots$, by 50 to 100 times up to a value which is, in every case, approximately equal or even higher than the critical density, at which the collectivization of the next electron level of the ion can occur.

Respectively, with a further increase in the drift motion energy (in the given case, at $E_K = 50 \text{ keV}$), the positions of minima shift to $n_{e(Z^*=2)} \approx 4 \times 10^{26} \text{ cm}^{-3}$, $n_{e(Z^*=3)} \approx 6 \times 10^{26} \text{ cm}^{-3}$, and $n_{e(Z^*=4)} \approx 10^{27} \text{ cm}^{-3}$. Each of

these values is considerably more than the destruction threshold of the next (deeper) electron state. For such an energy of the drift motion of electrons, there will occur the avalanche-like ionization of atoms in the target and the avalanche-like self-compression of the target substance in the volume of a propagating plasma layer (in the volume of a spherical wave “scanning” the target).

It is natural that the same effect will be observed for heavy atoms. In Figs. 11.20 and 11.21, we give the results of calculations of the total interaction energy of ions in a copper target with a degenerate electron gas moving across a target with the drift velocity corresponding to the drift motion energy E_K .

In Fig. 11.20, we consider the case of a relatively small energy E_K ($10 \text{ keV} \leq E_K \leq 100 \text{ keV}$), and, in Fig. 11.21, we present the results of analysis for a great interval ($0.05 \text{ MeV} \leq E_K \leq 3 \text{ MeV}$) of values of the kinetic energy of drift motion, E_K .

It follows from Fig. 11.20 that, with increase in the drift energy, there occurs a continuous shift of the position and depth of the minimum of the total interaction energy. The analysis of this figure indicates that, for nonrelativistic drift velocities, the equilibrium electron density n_e defined by the position of the energy minimum $U_{e\Sigma}/Z$ depends almost

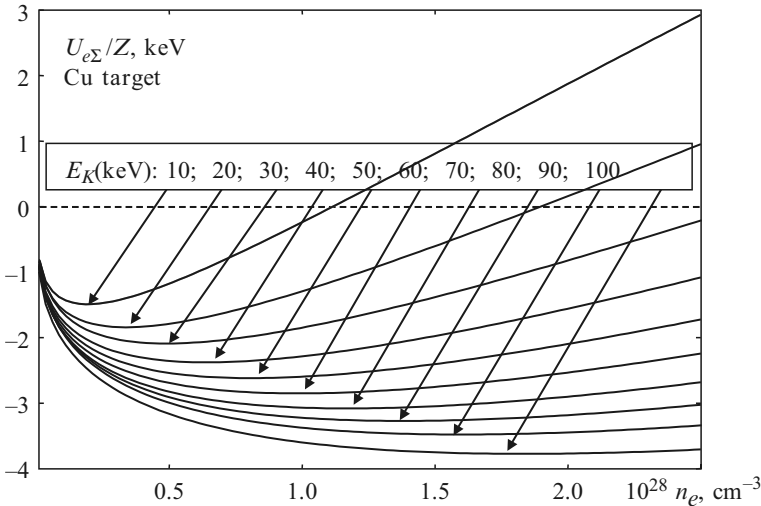


Fig. 11.20. The total energy of the interaction of ions of a copper target with the degenerate electron gas moving across the target with the drift velocity corresponding to the energy of drift motion E_K ($10 \text{ keV} < E_K < 100 \text{ keV}$).

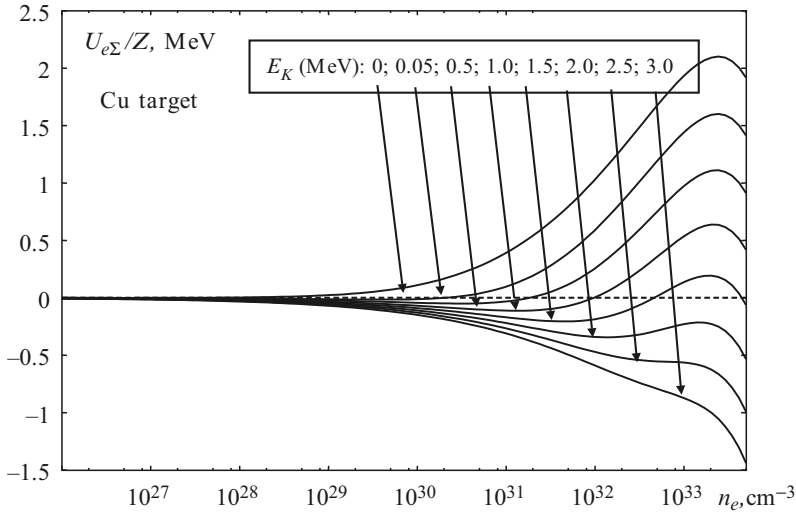


Fig. 11.21. The total energy of the interaction of ions of a copper target with the degenerate electron gas moving across the target with the drift velocity corresponding to the energy of drift motion E_K ($0.05 \text{ MeV} < E_K < 3 \text{ MeV}$).

linearly on the energy of drift motion E_K [and on the square of the drift velocity, $v^2 \equiv (dR/dt)^2$]:

$$n_e/E_K \approx \text{const} = \beta_Z. \quad (11.113)$$

This relation yields

$$n_e \approx (2\beta_Z/m_e)v^2 \equiv \alpha_Z(dR/dt)^2. \quad (11.114)$$

An analogous relation holds true for the other substances which are metals at such values of the electron density.

In particular, the approximation coefficient following from this calculation for Cu is $\alpha_{Cu} \approx 0.6 \times 10^8 \text{ s}^2/\text{cm}^5$.

The other important parameter is the depth of a potential well characterizing the total specific interaction energy $U_{e\Sigma}/Z$. It follows from the analysis performed that, for nonrelativistic drift energies, it increases by one order, being approximately proportional to the cubic root of the drift energy $|U_{e\Sigma}/Z| \sim (E_K)^{1/3}$, and is described by the approximation formula

$$U_{e\Sigma}/Z \approx -b(dR/dt)^{2/3}. \quad (11.115)$$

For copper targets, the approximation coefficient is $b \approx 3 \times 10^{-15} \text{ g cm}^{4/3}/\text{s}^{4/3}$.

The analysis of the specific binding energy $U_{e\Sigma}/Z$ for the relativistic region of the drift energy (see, e.g., Fig. 11.21 for $E_K > 0.5 \text{ keV}$) leads to

the formula $U_{e\Sigma}/Z \approx -kE_K$. In particular, for a copper target, the approximation coefficient is $k \approx 1/6$.

It follows from Fig. 11.21 that, in the region of very large drift energies (for Cu, at energies exceeding the threshold value $E_{K(\text{thresh})} \approx 2 \text{ MeV}$), the dependence of the binding energy on the density of electrons changes basically. If the energy of drift motion is less than the threshold value $E_{K(\text{thresh})}$, the general regularity of the dependence of the total energy on the density of the degenerate electron gas is characterized by the presence of the critical density of the electron gas, after the attainment of which the process of “falling of electrons onto a nucleus” is realized. But if the energy of drift motion exceeds the threshold value $E_{K(\text{thresh})}$, such a threshold is absent in the general case.

Consider these peculiarities in more details.

Let $E_K < E_{K(\text{thresh})}$. With increase in the density of a degenerate electron gas above the equilibrium electron density $n_{e(Z)}$, the energy of the system sharply grows at first. But, after the attainment of the threshold density $n_{ecr} \approx 10^{36} Z^{-2} \text{ cm}^{-3}$, it irreversibly decreases, which corresponds to the state of collapse and to the process of “falling of electrons onto a nucleus”. It is seen that, in this case, the region of a collapse is separated from the “normal” state of the electron-nucleus system by a high potential barrier, whose height decreases with increase in the energy of drift motion. This situation is presented, in particular, in Fig. 11.21.

At a very large energy of drift motion E_K of the degenerate electron gas [at $E_K > E_{K(\text{thresh})}$], the situation becomes basically different. In that case where the position of the minimum $n_{e(Z)}$ approaches the threshold density of the degenerate electron gas $n_{ecr} \approx 10^{36} Z^{-2} \text{ cm}^{-3}$ with increase in the energy of drift motion E_K , the behavior of the total interaction energy $W_{e\Sigma}$ is basically changed. Instead of the increase in $W_{e\Sigma}$ at $n_e \rightarrow n_{ecr}$, there occurs a very fast and unbounded decrease in $W_{e\Sigma}$. Such a decrease in the interaction energy $W_{e\Sigma}$ induces the irreversible collapse of the electron-nucleus system which begins, in fact, at a small density of the electron gas and is realized without any action of a potential barrier which would hamper or decelerate such a collapse. To ensure such a mode of the irreversible self-compression of a copper target, a single condition must be necessarily satisfied: the kinetic energy of the drift motion should exceed the threshold value $E_{K(\text{thresh})}$ [for a copper target, $E_{K(\text{thresh})} \approx 2 \text{ MeV}$]. The threshold kinetic energy of drift motion will be less or more than $E_K \approx 2 \text{ MeV}$ for targets manufactured from atoms, respectively, heavier or lighter than Cu. In particular, $E_{K(\text{thresh})} \approx 1 \text{ MeV}$ for a target containing U, and $E_{K(\text{thresh})} \approx 6 \text{ MeV}$ for a target containing Al.

Thus, the presence of the drift motion of electrons leads to that the process of collapse of an electron-nucleus system can begin at a small initial electron density (in the limit even at the ordinary “solid” density characteristic of undeformed metals) and then spontaneously reaches the final state of collapse “by passing,” in turn, all stages of the fast self-compression.

As was shown in Ref. 56 and considered above, such a Coulomb collapse of an electron-nucleus system induces the increase in the binding energy of this system and opens the possibility of the energy-gained synthesis of heavy and superheavy nuclei with $A > 300$. In the mentioned work, it is shown that the practically complete screening of the charge of nuclei in the extremely compressed relativistic degenerate gas of electrons suppresses the action of a Coulomb barrier and accelerates the reactions of synthesis of such nuclei.

11.2.4. Formation and the Motion of a Scanning Spherical Layer of a Degenerate Electron Gas in a Condensed Target

The discovered interconnection between the kinetic energy of drift motion and the possibility for an electron-nucleus system to collapse allows us to raise the question about whether the bounded 3D degenerate electron gas can be accelerated in a condensed medium up to such an energy of drift motion, when the conditions for the formation of a barrierless electron-nucleus collapse are satisfied.

From the purely formal viewpoint, the process of drift of conduction electrons can be realized, for example, at the expense of the application of a very strong electric field to a target and the generation of the electric current with extremely high density in it. Let us take the standard coupling of the density of an electric current j_e with the drift velocity v_d and the density of electrons n_e ,

$$j_e = en_e v_d. \quad (11.116)$$

We see that, in order to attain the drift velocity close to a relativistic one, it is necessary to ensure the current density on the level of $j_e \approx 10^{14}$ A/cm² at the initial density of conduction electrons in metals $n_e \approx 10^{23}$ cm⁻³. The absolute unreality of such force method for the “acceleration” of electrons in metals is obvious. Such values of j_e are unattainable, because any target is exploded already at the critical density of the electric current $j_{ecr} \approx 10^9$ A/cm² at the expense of the electron-plasma instability. The mechanism of such an instability consists in that, at such a critical current density, the velocity of the drift motion of conduction electrons turns out to be synchronized with the velocity of plasma oscillations, in which ions and

remaining electrons take participation. This corresponds to the Cherenkov instability accompanied by a coherent transfer of energy from electrons to the plasma and leads to the instantaneous heating and destruction of a target. Such an effect occurs, for example, upon the explosive emission with the typical critical density $j_{ecr} \approx 10^8$ to 10^9 A/cm².

It is obvious that the acceleration of electrons should be realized at the expense of the inner mechanisms of the mutually consistent motion of all degenerate electrons and nuclei. Only in such a system, the drift motion of the dense gas of degenerate electrons with a great Fermi energy will not cause their heating due to the Pauli principle.

Consider the peculiarities of a manifestation of the surface energy and the forces of surface tension on the boundary of the degenerate electron gas, whose volume contains the ordered ion lattice.

The main condition for the stable existence of a neutralized degenerate electron gas is the requirement of its equilibrium relative to the action of the Coulomb forces of the interaction of electrons and ions and the Fermi forces defining the kinetic pressure. This condition defines the requirement to minimize the volume occupied by the degenerate gas and is reduced to the condition of the equality of the pressures induced by these forces, ($|dU_{eQL}/dV_0| = |dU_{eF}/dV_0|$), on the surface of a Wigner-Seitz cell and on the external surface of the entire region which contains the gas of degenerate electrons.

The next important aspect is the account of forces defining a form of the surface of the region containing the degenerate gas. If the form of the volume containing this degenerate gas deviates from a sphere, its surface will undergo the action of a compressing pressure which is analogous to the surface tension and tries to decrease the surface. For this reason, the surface of a localized degenerate gas with high density always has the form close to a sphere.

One can determine this pressure by using Eq. 11.98 for the total binding energy $U_{e\Sigma}/Z$ modified with regard to the influence of the drift motion of electrons to the target center.

For a spherical layer with volume V , the total binding energy is $(U_{e\Sigma}/Z)Vn_e$. By using the above-deduced approximation dependence $U_{e\Sigma}/Z \approx -b(dR/dt)^{2/3}$ Eq. 11.115, we get the total pressure on the plasma layer surface as

$$P_\Sigma = U_{e\Sigma}n_i \approx -b(dR/dt)^{2/3}n_e. \quad (11.117)$$

There are two main results of the action of the surface tension force.

1. It tries to minimize the area of the region with the degenerate gas of electrons keeping the volume of this region to be constant. Since the

analogous pressure will act from all sides, the external and internal surfaces of this region try finally to take the spherical form.

2. Under the action of the same force, the thin spherical layer with a degenerate gas of electrons begins to decrease its area by moving to the target center.

Consider the result of the action of this force in the case where the gas of degenerate electrons occupies the volume in the form of a thin concentric spherical layer between two concentric spheres with radii R and $R - \Delta R$.

Each of the surfaces (inner and outer) undergoes the action of the identical pressure $\vec{P}$ directed normally to the surface inward the layer. Due to a difference in the areas of the inner and outer surfaces, the summary (resulting) force

$$F_{\Sigma} = -\{4\pi R^2 P_{\Sigma} - 4\pi(R - \Delta R)^2 P_{\Sigma}\} \approx -8\pi R \Delta R P_{\Sigma} \quad (11.118)$$

acts on the whole spherical layer and is directed to the center of symmetry of this layer, i.e., to the geometric center of the sphere. Under its action, the spherical layer with a high density of degenerate electrons with the total mass $M_{e\Sigma} = 4\pi R^2 \Delta R m_e n_e$ will move to the center.

The equation describing the radial nonrelativistic motion of the spherical layer is as follows:

$$d^2 R/dt^2 = F_{\Sigma}/M_{e\Sigma} = -2b(dR/dt)^{2/3}/Rm_e. \quad (11.119)$$

This nonlinear equation is true, as long as the velocity of the layer undergoing compression remains nonrelativistic.

By multiplying both sides of Eq. 11.119 by $(dR/dt)^{1/3}$, we get

$$d\{(dR/dt)^{4/3}\}/dt = -(8b/3m_e)(dR/dt)/R \quad (11.120)$$

and now can easily obtain the dependence of the running motion velocity dR/dt of the plasma layer undergoing compression on its radius R .

With regard to the initial conditions $dR/dt|_{t=0} = v_0$ and $R|_{t=0} = R_0$, relation Eq. 11.120 yields

$$dR/dt = -\left\{v_0^{4/3} + (8b/3m_e) \ln(R_0/R)\right\}^{3/4}. \quad (11.121)$$

For a target manufactured from Cu, we get

$$dR/dt \approx -\left\{v_0^{4/3} + 10^{13} \ln(R_0/R)\right\}^{3/4} \text{ cm/s}. \quad (11.122)$$

It is seen that, at $R_0/R > 2.7$, the radial velocity of the plasma layer (even at $v_0 = 0$) approaches the velocity of light (in this case, $|dR/dt| >$

6×10^9 cm/s), and a correct description of the process of compression requires to use a relativistic equation of motion.

Consider the other limiting case corresponding to the movement of a plasma layer undergoing compression which has a very high density of electrons (and, respectively, a high local density of ions). Such a mode of compression corresponds to the relativistic gas of degenerate electrons.

The relativistic equation of motion for this layer has the form

$$dp_\Sigma/dt = F_\Sigma \approx -8\pi R\Delta R P_\Sigma, v \approx c. \quad (11.123)$$

Here,

$$p_\Sigma = \gamma M_{e\Sigma} v \approx 4\pi\gamma R^2 \Delta R m_e n_e c \quad (11.124)$$

is the total relativistic momentum of the layer undergoing compression, γ is Lorentz-factor of moving relativistic electrons, $P_\Sigma \approx (U_{e\Sigma}/Z)n_e \approx -kE_K n_e$ is the surface pressure of the relativistic degenerate gas, and $E_K = \gamma m_e c^2$ is the total drift energy motion of every electron.

With regard to the fact that the velocity of relativistic particles with $\gamma \gg 1$ is almost constant and is equal to c , the equation of motion for the plasma layer reduces to the equation for the Lorentz-factor.

In view of the relation $R \approx R_0 - ct$, the solutions of the equation of motion

$$d\gamma/dt \approx -2kE_K/Rm_e c = -2k\gamma c/(R_0 - ct) \quad (11.125)$$

with regard to the initial conditions $[\gamma(t = t_0) = \gamma_0 \gg 1, R(t = t_0) = (R_0 - ct_0) \text{ at } t = t_0]$ take the form

$$\begin{aligned} \gamma(t) &\approx \gamma(t_0)[(R_0 - ct_0)/(R_0 - ct)]^{2k} \\ &\approx \gamma(t_0)[(R_0 - ct_0)/(R_0 - ct)]^{1/3}, \end{aligned} \quad (11.126)$$

$$E_K(t) = \gamma(t)m_e c^2. \quad (11.127)$$

As soon as the energy of linear (drift) motion of the spherical plasma layer will reach the threshold value E_K , the irreversible collapse of the electron-nucleus system will occur.

It is necessary to note that, besides the force of surface tension, the other forces are present which can efficiently compress and accelerate the plasma layer.

For example, if the process of compression of the plasma layer is developing simultaneously with the passage of the electric current along a target, the plasma layer will be additionally compressed and accelerated at the expense of the pressure of the magnetic field. Such a situation will be

realized in the case where the primary near-surface plasma wave is formed as a result of the action of a shock wave induced, for example, by the deceleration of a symmetric electron beam on the anode. In this case, the electric current J begins to run along the anode, and its magnetic field creates the pressure

$$P_J = H_\varphi^2/8\pi = J^2/2\pi R^2 c^2 \quad (11.128)$$

on the plasma layer surface.

This pressure acts on the part of the spherical layer surface maximally remote from the current axis, which reflects the cylindrical geometry of the passage of the electric current. At the first glance, it seems that such a pressure will accelerate only this part of the surface. However, it is easy to verify that, in the case where the layer consists of a strongly degenerate electron gas, there will occur not only the acceleration of this local part of the layer and the deformation of the whole layer related to this acceleration, but the symmetric uniform compression and the acceleration of the layer without any change in the form of its surface. This is conditioned by the action of the very strong surface tension, hampering the increase in the surface area, in the bounded volume containing both the gas of degenerate electrons and ions neutralizing the electrons. If the additional external pressure is less than the pressure induced by the forces of surface tension, then any symmetric (in the counterdirections along any of the diameters of the spherical layer) action on small parts of the spherical layer with the degenerate gas will lead to the uniform compression and the acceleration of the whole layer.

One more mechanism of compression and acceleration of the plasma layer can be related to the action of a light-related reactive or ablation-induced pressure. This pressure arises upon the emission of hard quanta or other nuclear particles with energy $U_{e\Sigma}/Z$ emitted in the process of release of the binding energy upon the compression of the plasma layer. This pressure can exceed the pressure caused by the above-considered forces of surface tension in the bounded neutralized plasma layer.

As one more (and, possibly, main) reason for the accelerated motion of the plasma layer, we indicate the nuclear reactions running in the volume of this compressed layer. The binding energy released in these reactions can also be a reason for the additional acceleration of this layer. This mechanism of acceleration can become principal, if we take into account the circumstance that many of the alternative channels of the energy release in such extremely symmetrized system can be strongly suppressed. In the first turn, this concerns the ban on the emission of hard γ -quanta upon monopole transitions. These questions were partly considered in Sec. 11.1.6

Thus, the results defined by Eqs. 11.121, 11.122, and 11.127 characterize only the minimum acceleration rate which can be, in fact, significantly higher.

11.2.5. Motion of the Ion (Nuclear) Component of a Target in a Scanning Spherical Layer

Adaptive Reaction of the Ion Component of a Target on the Motion of a Scanning Plasma Layer with High Density. Consider the question about the mutual influence of ions of a target and a rapidly moving spherical layer of degenerate electrons. This question is exclusively important for both the comprehension of the collapse formation dynamics and the analysis of those nuclear processes which are running in the compressed medium. Consider several alternative approaches to this problem (which involve plasma oscillations and shock waves).

In the first turn, we consider the basic aspect related to the specificity of the influence of the fast motion of the layer with a high density of the degenerate electron gas on a state of the ion (nuclear) system in a condensed target. As shown above, the equilibrium ion density grows sharply in a moving plasma layer undergoing compression. At the first glance, it seems that the expenditures of energy on the relevant shift of ions to a new equilibrium position must lead to the deceleration of a propagating electron layer. However, it is easy to verify that the work on the displacement of ions is actually performed, in this case, by ions themselves, rather than by electrons. Indeed, because the “superfluous” collectivized electrons belonging to the composition of the degenerate gas in the case where the density of the gas of these electrons is higher than the equilibrium density for a given configuration of ions are simultaneously attracted by different ions from different sides, this combination of interactions is equivalent to the mutual attraction of these ions without essential change in the state of electrons (i.e., without essential decrease in the drift velocity). Such a circumstance means that the increase in the ion density occurs at the expense of the indirect interaction of the very ions, though upon the direct participation of electrons.

Let us estimate the parameters of the process of adaptation of ions to a variation in the electron density. According to the model under consideration, the target prior to the “arrival” of a spherical plasma layer undergoing compression was an ordinary metal with the ion density $n_{i0} \approx 10^{23} \text{ cm}^{-3}$. We note that, if the target would be manufactured of a dielectric or a semiconductor, the subsequent scenario of its evolution would be approximately the same. To provide the continuous acceleration of the plasma layer and to enhance its density, it is necessary that the ions in the volume of this layer have time, for the time interval of its passage, to be ionized and to

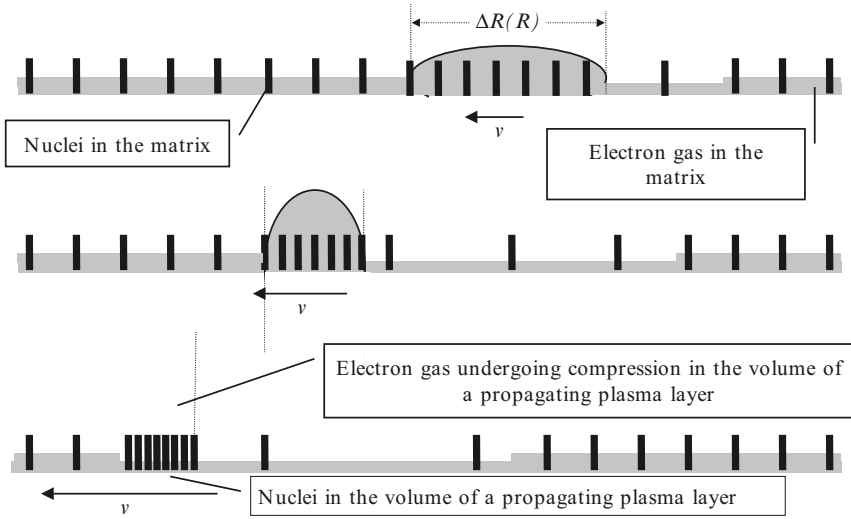


Fig. 11.22. Successive stages of the compression of a spherical plasma layer propagating to the center of a target.

be “squeezed” up to the equilibrium ion density $n_i = n_e/Z$ which corresponds to the running electron density n_e in the volume of the plasma layer. Such a process of compression of the spherical plasma layer is symbolically presented in Fig. 11.22.

At first, we make the qualitative estimations justifying the possibility of both an adequate fast “reaction” of the ion component of the target on the motion of electrons and the increase in their concentration at any specific place of the target. To this end, we will determine how much time the transition of ions to the equilibrium position upon a change in the electron density takes.

In a neutral plasma, the duration Δt_1 of a local displacement of any ion from a nonequilibrium position up to the equilibrium position is independent of the displacement amplitude and is defined by a quarter of the period of plasma oscillations $T_i = 1/\omega_{i0}$ of the ion component of the compressed plasma layer, whose frequency $\omega_{i0} = (4\pi n_i Q_i^2/M_i)^{1/2}$.

Here, $n_i = n_e/Z$, $Q_i = Ze$, $M_i = AM_0$. Starting from this condition, we get

$$\Delta t_1 \approx 1/4\omega_{i0} = [(A/Z)M_0/64\pi n_e e^2]^{1/2} \approx (M_0/30\pi n_e e^2)^{1/2}. \quad (11.129)$$

The minimum kinetic energy W_{KZ} of every nucleus in the compressed electron-nucleus plasma corresponds to the energy of zero oscillations of a 3D oscillator

$$W_{KZ(\min)} = 3\hbar\omega_{i0}/2 \approx 3\hbar e(\pi n_e Z/2M_0)^{1/2}. \quad (11.130)$$

At a low temperature (at $T \ll T_F = Z^2 e^2 / 2L_0$), the system of nuclei in the degenerate gas takes the form of a cubic lattice with period $L_0 = 2(Z/n_e)^{1/3}$. With increase in temperature, the amplitude of oscillations grows, and, at $T > T_F$, the lattice melts and transforms into a nuclear liquid being positioned in the neutralizing compressed electron gas. This conclusion follows directly from the Lindemann criterion (see Ref. 62), according to which the melting temperature of a lattice T_F is determined from the condition that every nucleus is displaced by a quarter of the period. Under such a displacement, $Z_{\text{eff}} = Z/2$, $L = L_0/2$. At $Z = 29$ and $n_e \approx 10^{30} \text{ cm}^{-3}$, the melting temperature of the nuclear lattice is sufficiently high ($T_F \approx 100 \text{ keV}$) and increases with Z and n_e .

Eqs. 11.129 and 11.130 yield that the duration of a local displacement of a nucleus up to the equilibrium position is at most

$$\Delta t_1 \approx 3 \times 10^{-17} \text{ s}, W_{KZ(\text{min})} \approx 20 \text{ eV at } n_e \approx 10^{26} \text{ cm}^{-3}, \quad (11.131)$$

$$\Delta t_1 \approx 3 \times 10^{-19} \text{ s } W_{KZ(\text{min})} \approx 2 \text{ keV at } n_e \approx 10^{30} \text{ cm}^{-3}, \quad (11.132)$$

$$\Delta t_1 \approx 3 \times 10^{-21} \text{ s}, W_{KZ(\text{min})} \approx 0.2 \text{ MeV at } n_e \approx 10^{34} \text{ cm}^{-3}. \quad (11.133)$$

Let us compare these durations of a displacement of a nucleus with the time

$$\Delta t_2 \approx \Delta R/v \approx \Delta R_0(n_{e0}/n_e)^{1/3}/v \quad (11.134)$$

of the “passage” of a plasma layer undergoing compression through the place where a specific ion is located (in this case, $v \approx c$ at $n_e \geq 10^{30} \text{ cm}^{-3}$). The layer thickness $\Delta R \approx \Delta R_0(n_{e0}/n_e)^{1/3}$ is defined by the initial thickness of the layer ΔR_0 (at $R = R_0$), the initial density $n_{e0} \approx 6 \times 10^{22} \text{ cm}^{-3}$, and the running density n_e of electrons in the layer.

Of course, the value of Δt_1 defines the minimum duration of a rearrangement of the ion (nuclear) system only if there exists a mechanism of consistent redistribution of the kinetic energy of ions in the process of movement to a position of local equilibrium. This consistent movement can be related to various processes. The most grounded consideration corresponds to a collective motion of nuclei of the type of a shock wave which will be analyzed below.

Let us take that the initial (at $R = R_0$) thickness of a spherical plasma layer is $\Delta R_0 \approx 10 \mu\text{m}$. Then, at $n_e \approx 10^{26} \text{ cm}^{-3}$, we have $\Delta R \approx 1 \mu\text{m}$ and $\Delta t_2 \geq 3 \times 10^{-15} \text{ s}$. If $n_e \approx 10^{30} \text{ cm}^{-3}$, then $\Delta R \approx 400 \rho\text{A}$ and $\Delta t_2 \approx 10^{-16} \text{ s}$; and if $n_e \approx 10^{34} \text{ cm}^{-3}$, $\Delta R \approx 4 \rho\text{A}$ and $\Delta t_2 \approx 10^{-18} \text{ s}$.

Respectively, at the initial thickness of the layer $\Delta R_0 \approx 1 \mu\text{m}$, we get $\Delta R \approx 0.1 \mu\text{m}$ and $\Delta t_2 \geq 3 \times 10^{-16} \text{ s}$ at $n_e \approx 10^{26} \text{ cm}^{-3}$, $\Delta R \approx 40 \rho\text{A}$ and $\Delta t_2 \approx 10^{-17} \text{ s}$ at $n_e \approx 10^{30} \text{ cm}^{-3}$, and $\Delta R \approx 0.4 \rho\text{A}$ and $\Delta t_2 \approx 10^{-19} \text{ s}$ at $n_e \approx 10^{34} \text{ cm}^{-3}$.

It follows from these estimations that $\Delta t_2 \gg \Delta t_1$. This means that the time Δt_2 of the passage of a spherical plasma layer undergoing compression through the place, where any nucleus is positioned in the matrix, turns out to be sufficiently large to ensure a displacement of this nucleus for the time Δt_1 to the new position which corresponds to the equilibrium of the superdense medium with ion density $n_i = n_e/Z$ which is created for a short time in the volume of a propagating self-compressing spherical layer. It is seen that the ion system is able to adequately and rapidly respond to the drift motion of electrons.

Formation and Movement of a Shock Wave Synchronized with the Motion of a Scanning Plasma Layer. In the case where the initial state of the ion (nuclear) subsystem corresponds to the unperturbed (undeformed) matrix of a target, the above-considered mechanism of adaptive displacement of ions describes quite satisfactorily their response to the accelerated motion of electrons in the plasma layer volume.

A basically different scenario should be considered in the case where the process of formation of the region of primary ionization on the target periphery is accompanied by a simultaneous impact (pulse) action on the ion system. This is the most real case.

In view of both the very strong interaction of electrons and nuclei at their great concentration and the nonstationarity of the system under study, it is obvious that, in this case, the adequate description of the process of self-compression of the target substance in the volume of a spherical scanning layer should be based on a formalism considering the collective and self-consistent motion of electrons and nuclei.

The accelerated motion of the electron layer being in the state of a compressed degenerate electron gas is accompanied by the density wave of ions which is, in essence, an electron-nucleus shock wave propagating in this electron gas.

First, we will determine the parameters of an ordinary (not shock) wave in a compressed degenerate gas.

It is well known that the velocity v_{L1} of the longitudinal sound in a gas is defined by its volumetric compressibility $V(\partial^2 U_{eF}/\partial V^2)$ and the density of the medium $\rho = M_i n_e/Z$. For a nonrelativistic degenerate gas, this velocity is

$$\begin{aligned} v_{L1} &= [V(\partial^2 U_{eF}/\partial V^2)/\rho]^{1/2} = (3\pi^2 n_e)^{1/3} \hbar (Z/3M_i m_e)^{1/2} \\ &\approx 0.02 n_e^{1/3} \text{ cm/s.} \end{aligned} \quad (11.135)$$

Such a velocity corresponds, in particular, to the region of a compressed electron gas before the leading edge of a shock wave. With increase

in the density n_e , the velocity v_{L1} grows and reaches the value $v_{L1} \approx 2 \times 10^8$ cm/s for the limiting density of a nonrelativistic gas (at $n_e \approx 10^{30}$ cm $^{-3}$).

For a relativistic degenerate gas of electrons (at $n_e \geq 10^{31}$ cm $^{-3}$), the sound velocity is

$$v_{L1} = (3\pi^2 n_e)^{1/6} (Z\hbar c / 3M_i)^{1/2} \approx 2.5 \times 10^3 n_e^{1/6} \text{ cm/s}. \quad (11.136)$$

It follows from Eq. 11.136 that the sound velocity in a relativistic degenerate gas of electrons is $v_{L1} \approx 4 \times 10^8$ cm/s at $n_e \approx 10^{31}$ cm $^{-3}$ and reaches its limiting value $v_{L1} \approx 5.4 \times 10^9$ cm/s at $n_e \approx 10^{38}$ cm $^{-3}$.

We note that the account of the degeneration of the nuclear component of a plasma in the extremely dense relativistic medium leads to a greater limiting sound velocity $v_{L1} = c/\sqrt{3} \approx 1.7 \times 10^{10}$ cm/s (e.g., see Chapter 15 in the E. Teller's book [Ref. 39]).

In this compressed gas, one can observe the propagation of shock waves. Such a wave is formed from the initial perturbation which has acted on the target surface and, along with other things, has led to the formation of the initial region of ionization in the target. Let us take into account that the shock wave velocity D can significantly exceed the sound velocity v_{L1} . Then, at a great density of the degenerate electron gas, a subrelativistic or relativistic shock wave of ions (nuclei), which is synchronous with the motion of the electron layer, can be formed.

It is obvious from the qualitative considerations that a shock wave cannot have the mean radial velocity exceeding the drift velocity of the motion of a bounded (along the radius) spherical layer of the compressed electron gas. This conclusion follows at once from the evident fact that a large velocity D of a shock wave can be reached only upon its propagation in the compressed electron gas. It is clear that as soon as the leading edge of a shock wave attain the forward front of the layer undergoing compression (i.e., its inner surface nearest to the target center), then the shock wave goes at once into the uncompressed medium being in front of this layer, where the shock wave velocity abruptly decreases. After such a deceleration, the shock wave "returns" into the volume of the plasma layer moving in the same direction. However, having "returned," it begins again to accelerate and again overruns the movement of the spherical layer of the compressed electron gas. Then the process is repeated.

Such qualitative consideration shows that the leading edge of a shock wave turns out to be localized near the forward front of a propagating spherical layer of the compressed electron gas.

Now we consider the quantitative parameters of a shock wave which is "developed" against the background of the degenerate electron gas undergoing compression.

The process of development of any shock wave is characterized by several basic relations (see, e.g., Ref. 39). Among them, we mention the shock adiabat

$$P_2/P_1 = [(\gamma + 1)V_1 - (\gamma - 1)V_2]/[(\gamma + 1)V_2 - (\gamma - 1)V_1], \quad (11.137)$$

following from the standard Rankine–Hugoniot relations, the equation of state for a substance, and the equations of state for the temperature,

$$T_2/T_1 = [2\gamma M^2 - (\gamma - 1)][(\gamma - 1)M^2 + 2]/(\gamma + 1)^2 M^2, \quad (11.138)$$

pressure,

$$P_2/P_1 = 2\gamma M^2/(\gamma + 1) - (\gamma - 1)/(\gamma + 1), \quad (11.139)$$

and density

$$\rho_2/\rho_1 = n_{e2}/n_{e1} = (\gamma + 1)M^2/[(\gamma - 1)M^2 + 2]. \quad (11.140)$$

These equations define unambiguously the connection of the state of the substance before the leading edge of a shock wave (index 1) and after it (index 2).

In the case where the phase states of the substance before the leading edge of a shock wave and after it are the same (like in the case under consideration, when nuclear reactions do not begin else and the densities of the degenerate electron gas are essentially different on both sides of this front), the equation of shock adiabat reads

$$P_2/P_1 = [(\gamma + 1)/n_{e1} - (\gamma - 1)/n_{e2}]/[(\gamma + 1)/n_{e2} - (\gamma - 1)/n_{e1}]. \quad (11.141)$$

Here, γ is the parameter of the Poisson adiabat $p \sim 1/V^\gamma$ (or $p \sim n_e^\gamma$); P_2 and P_1 are the pressures, respectively, behind and before the leading edge of a shock wave; $V_2 \sim 1/n_{e2}$ and $V_1 \sim 1/n_{e1}$ are, respectively, the specific volumes behind and before the leading edge of a shock wave; $M = D/v_{L1}$ is the Mach number depending on the shock wave velocity D and the sound velocity v_{L1} before it.

As shown in Sec. 11.2.3, the equation of state of a degenerate gas in the presence of a drift motion corresponds to the presence of a local energy minimum $U_{e\Sigma}/Z$ (as a function of the electron gas density $n_e = Z/V$) which shifts to the region of a higher electron density with increase in the drift energy E_K . This result is presented in Figs. 11.19–11.21 and 11.23.

In view of the fact that the derivative of the energy $U_{e\Sigma}(n_e)/Z$ with respect to the density n_e of the electron gas defines the dependence of the pressure of this gas on its density, namely

$$P(n_e) = -d\{U_{e\Sigma}(n_e)/Z\}/dV = (1/V)^2 d\{U_{e\Sigma}(n_e)/Z\}/dn_e, \quad (11.142)$$

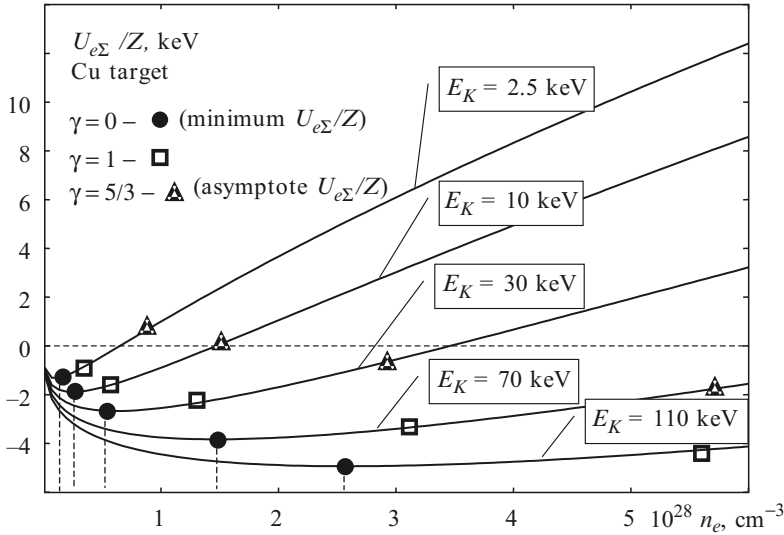


Fig. 11.23. Energy $U_{e\Sigma}/Z$ and the adiabat exponent γ of a degenerate nonrelativistic electron gas in the volume of a Wigner–Seitz cell Versus the concentration of electrons n_e for a specific energy of drift motion E_K . The minimum energy $U_{e\Sigma}/Z$ corresponds to both the concentration $n_e = n_e(Z)$ and $\gamma = 0$; the critical value $\gamma = 1$ corresponds to $n_e(\gamma = 1) > n_e(Z)$; and the asymptotic value of γ for a nonrelativistic gas is $\gamma_{Fnr} = 5/3$ and corresponds to the concentration of electrons $n_{e(\gamma=5/3)} > n_{e(\gamma=1)} > n_{e(\gamma=0)}$.

it becomes obvious that the position of the energy minimum as a function of the density n_e at $n_e = n_e(Z)$ corresponds to the zero pressure $P(n_e Z) = 0$, and the gas pressure at $n_e \gg n_e(Z)$ corresponds to the limiting Fermi pressure of a degenerate nonrelativistic gas $P(n_e) \sim n_e^{5/3}$ (this question was considered in Sec. 11.1.5).

It follows from the formal dependence of the Poisson adiabat $P(n_e) \sim n_e^\gamma$ on the parameter γ that the energy minimum $U_{e\Sigma}(n_e)/Z$ at $n_e = n_e(Z)$ corresponds to $\gamma = 0$. Respectively, at $n_e \gg n_e(Z)$, $\gamma \rightarrow 5/3$ in a nonrelativistic gas.

These circumstances allow us to conclude that, at any energy of drift motion, the adiabat parameter γ depends on the electron gas density and can take all intermediate values in the interval $0 \leq \gamma \leq 5/3$, including $\gamma = 1$. As known, the last value plays a particular role in the physics of shock waves. For a completely ionized copper target, these results are presented in Fig. 11.23.

We note that, from the formal viewpoint, γ at $n_e < n_e(Z)$ is a complex-valued number. This circumstance reflects those limitations ($\gamma > 0$, $\text{Im}\gamma = 0$) which are always used implicitly upon the derivation of Eqs. 11.137–11.140 describing a shock wave.

The basic difference of these results from those in the case of an immovable degenerate gas (without drift) consists in the availability of only one fixed dependence of the gas energy $U_{e\Sigma}(n_e)/Z$ on its density. In such driftless gas, there exists the interval of values of the density n_e , for which the Poisson adiabat parameter is in the region near $\gamma \approx 1$. But this interval is fixed for a specific charge of ions and corresponds to a relatively low (on the scale of the processes of shock compression) density n_e which is near the equilibrium value n_{eZ} under the metallic type of the coupling of ions in the lattice. These results are presented in Figs. 11.15 and 11.16. Let the gas pressure grow (e.g., in the process of compression upon the propagation of a shock wave). Then, upon moving away from the steady-equilibrium value n_{eZ} to the side of large values of n_e , the adiabat parameter γ increases rapidly and becomes constant already at a small excess of the electron gas density n_e above n_{eZ} (for the electron gas, γ tends to $5/3$ or $4/3$), which restricts basically the degree of compression on the leading edge of a shock wave in a medium without the drift of electrons.

The same result follows from Eq. 11.140. For an intense shock wave with a high Mach number $M \gg 1$, the limiting increment of the density on the leading edge of a shock wave is defined by the value

$$\rho_2/\rho_1 = n_{e2}/n_{e1} = (\gamma + 1)/(\gamma - 1). \quad (11.143)$$

This relation is justified if $(\gamma - 1)M^2 \gg 2$.

In the case of the “ordinary” nonrelativistic or relativistic degenerate electron gas without drift, $\gamma_{\max} = 5/3$ or $\gamma_{\max} = 4/3$, respectively. This yields the well-known conclusion that the limiting compression of a substance in a shock wave is bounded by small values:

$$(\rho_2/\rho_1)_{\max} = (n_{e2}/n_{e1})_{\max} = 4 \text{ at } \gamma = 5/3 \quad (11.144)$$

$$(\rho_2/\rho_1)_{\max} = (n_{e2}/n_{e1})_{\max} = 7 \text{ at } \gamma = 4/3. \quad (11.145)$$

On the contrary, the increase in the electron gas density in the case under consideration leads to a growth of the velocity of its drift. This causes the shift of the energy minimum $U_{e\Sigma}(n_e)/Z$ as a function of n_e and, hence, to the shift of the region of values of n_e , where $\gamma < \gamma_{\max}$. It is very important that, in this case, the interval of values of n_e , for which $\gamma \approx 1$, shifts as well.

Now we consider the dynamics of the process of development of a shock wave. If the velocity of the leading edge of a shock wave has exceeded the drift velocity of electrons, then the effective value of the parameter γ begins to increase (in the limit, to $\gamma_{\max}$), which leads to a decrease in both the density on the leading edge of the shock wave and its velocity. However, upon such a deceleration relative to the drift velocity, the effective value

of the parameter γ begins to decrease. In accordance with Eq. 11.143, this leads at once to the sharp increase in the medium density on the leading edge of the shock wave:

$$(\rho_2/\rho_1)_{\max} = (n_{e2}/n_{e1})_{\max} \gg 1 \text{ at } \gamma \rightarrow 1. \quad (11.146)$$

This increase in the density leads, in turn, to an increase in both the shock wave velocity and the effective value of the parameter γ , which induces finally both a decrease in the density and the deceleration of the shock wave. Then the whole process is repeated.

It is seen that the process under consideration is self-controlled, is characterized by a negative feedback, and induces the synchronization of the drift velocity of the electron-nuclear plasma layer and the velocity of the leading edge of the shock wave.

It is easy to verify that such a result is not limited by the requirement $M \gg 1$. Indeed, relation Eq. 11.140 yields that, at $\gamma < 1$ and any $1 \leq M < \infty$, the condition for the supercompression on the leading edge of a shock wave,

$$\gamma \approx 1 - 2/M^2, \quad (11.147)$$

is always satisfied. In this case,

$$\rho_2/\rho_1 = n_{e2}/n_{e1} \rightarrow \infty. \quad (11.148)$$

Of interest is the circumstance that the temperatures of a compressed electron gas before the leading edge of a shock wave (T_1) and behind it (T_2) in the region of synchronized motion, where $\gamma \approx 1$, turn out to be approximately the same irrespective of the Mach number M . This result follows directly from Eq. 11.138 at $\gamma \approx 1$. It is natural that such a condition will hold only for that section of “acceleration” of a shock wave, where nuclear transformations accompanied by the energy release have not begun else.

Because the effective “thickness” of the leading edge of a shock wave is much less than the radius of a propagating spherical plasma layer, in the volume of which the shock wave is localized, it is quite rightful to use results (Eqs. 11.137–11.141, 11.143, 11.146, and 11.147) derived for a plane shock wave in the case of a spherical shock wave.

With regard to the character of the variation in the adiabat parameter γ , the above-mentioned circumstances lead at once to the conclusion on the possibility to realize the efficient process of formation of a spherical shock wave dynamically synchronized with the propagating density wave in a degenerate electron gas.

We note that the above-performed analysis is based on the use of the standard equations (Eqs. 11.137–11.141) describing a shock wave which are deduced for the nonrelativistic or slightly relativistic systems.

There are all the foundations to think that the qualitative character of both the development of a shock wave and its interaction with the propagating spherical plasma layer with extremely high electron density remains the same (i.e., the synchronization, mutual amplification, and a fast growth in the density) also in the case of relativistic shock waves (the latter correspond to the final stage of the evolution of the zone of a propagating collapse). These conclusions are based, in particular, on the results of the previous works on the relativistic hydrodynamics of shock waves in gravitating stars.

In view of the fact that the limiting velocity of sound in a superrelativistic gas at a large kinetic energy of nuclei is $v_{L1} = c/\sqrt{3} \approx 1.7 \times 10^{10}$ cm/s, it is obvious that the velocity of a shock wave, being always more than the sound velocity in the same medium, becomes very close to the light velocity. Thus, a shock wave in the volume of the layer becomes relativistic but cannot overrun the plasma layer squeezing to the center (by overrunning this layer, the wave goes into the uncompressed medium, where its velocity drops sharply, and the wave “returns” again into the layer volume). It is clear that the above-presented arguments on the synchronization of the motion of both the layer with an extremely high density of electrons and the shock wave are equally convincing in both the relativistic and nonrelativistic cases.

11.2.6. Regularities of the Scanning Synthesis and Peculiarities of the Products of a Collapse

General Regularities and a Mechanism of Scanning Synthesis upon a Pulse Action. The above-considered peculiarities of the motion of a thin closed plasma layer with a low initial density in the volume of a condensed target allow us to predict the probable scenario of the process of collapse and to advance a mechanism of the efficient nucleosynthesis. As a result of the manifestation of the latter, we indicate the “rearrangement” of about 10^{20} to 10^{21} nucleons in a single experiment. At first, these nucleons entered the nuclei of a target, but, after the completion of the process of collapse, they turned out to be arranged into groups in the form of newly synthesized isotopes.

The reason for the formation of the above-mentioned layer is the impact pulse symmetric correlated action on the surface of a target. Under

the action of this perturbation, a shock wave is formed in the near-surface layer of the target, and the thin layer of a partially ionized substance appears.

One of the main reasons for the motion of the layer is a very large force of surface tension acting on the bounded region of the neutralized degenerate electron gas. Under the action of this force, the outer and inner surfaces of the closed plasma layer acquire rapidly the form of ideal spheres. The same force compels this layer to squeeze with a high velocity to the target center. Simultaneously, the layer ionizes the target substance inside itself, and its density rapidly grows.

The interrelation between the internal energy of the plasma layer, its centripetal acceleration, and the force of surface tension leads to the synchronous self-compression, increase in the density of the degenerate electron gas in the volume of this layer, and formation of a self-controlled spherical wave with extremely high electron and ion densities (the “wave-shell”). This thin wave converging to the center scans all the target volume layerwise, by beginning from the region of formation of the wave and terminating in the region at the target center, where it undergoes a spherical collapse. In the process of compression of the ion (nuclear) component of the plasma layer, the decisive role is played by the shock wave which propagates synchronously with the motion of the plasma layer and is localized near its forward front. Upon the attainment of the critical kinetic energy E_K of the centripetal motion of the plasma layer, the Coulomb collapse of the electron-nucleus system occurs in its volume. One of the main conditions for a realization of such a collapse is the presence of a fast drift motion of the gas of degenerate electrons. The threshold value of the critical energy of such a motion E_K which is necessary for the realization of a collapse can be significantly reduced if we take into account the action of other mechanisms of compression of the propagating plasma layer (including the pressure of the magnetic field).

The results of calculations presented above indicate that the linear velocity of the motion of the spherical plasma layer can be close to the light velocity at the last stage of the compression.

At this stage, the density of the electron component of a self-controlled spherical wave with extremely high density corresponds to a relativistic degenerate gas, which allows realizing the state of Coulomb collapse of electrons and nuclei, screening the charges of nuclei in the compressed electron Fermi-condensate, and, finally, providing the efficient nucleosynthesis. If the density of the degenerate gas of electrons exceeds the threshold value depending on the charge of nuclei, then the fast synthesis of any nuclei (including heavy and anomalous nuclei) becomes possible in the scope of the “wave-shell”.

Peculiarities of the interaction of the scanning plasma layer with atoms of a target are defined by a number of successive processes:

- full ionization of the medium on the leading edge of the moving layer
- compression of the nuclear component in the volume of this layer up to the equilibrium density, which leads to the charge neutralization of the electron component and to the formation of the state of electron-nucleus collapse with the release of the binding energy
- cascade of nuclear transformations in the collapse zone
- decay of this collapse at the trailing edge of the plasma layer
- restoration of the condensed state of a target substance subjected to a successive chain of such critical actions

A part of the binding energy of the electron-nucleus system which is released in the process of formation of a collapse can lead to an additional acceleration of the scanning plasma layer. In the process of formation of the collapse state, there occurs the whole complex of nuclear transformations. Then their products remain near the place of their creation and correspond to a local position of the collapsing plasma layer at the moment of the realization of these nuclear transformations. The decay of the compressed state of the electron-nucleus plasma and the relaxation of this state to the “normal” one, i.e., to the uncompressed state of a target after the passage of a scanning wave of the superdense gas of electrons, is the endothermal process and causes the cooling and the localization of a part of the heavy products of nuclear transformations. It is natural that the element and isotope compositions of the target after the passage of the scanning plasma layer can very strongly differ from the initial composition prior to the passage of this layer.

The analysis of mechanical damages of the target after the completion of the process of collapse shows that, on the last stage of the motion of a collapsing spherical wave converging to the center of a target, there occurs the full mechanical cleavage of some substance from the remaining part of the target, being closer to the target surface. Due to a sufficiently high temperature, this concentric slit will be filled with the plasma possessing a density considerably less than that of a solid.

In the central part of the target, there occur the impact collision and the fusion (squeezing) of parts of the scanning plasma layer, which leads to both the formation of the zone of a spherical collapse and the attainment of the limiting compression at the center. The compression degree increases at the expense of the action of inertial forces.

The dynamics of processes inside this region is very complicated and requires an additional comprehensive study with regard to the processes of creation and annihilation of particles and the running of all kinds of nuclear

transformations (including, possibly, those with the creation of a quark-gluon plasma). It is obvious that, in this region, the most essential nuclear transformations can occur in this region, including the synthesis of a great number of superheavy nuclei and nuclei with anomalous isotope ratios, and a great amount of energy (possibly, in the form of hard γ -quanta, various kinds of neutrinos, and fast electrons and protons) can be released. After the termination of the action of inertial forces due to the heating at the expense of the running reactions, there occur the irreversible destruction of the zone of a spherical collapse and the ejection of the products of nuclear transformations into the environment. In any case, the lifetime of the state of collapse at the center of a target is considerably less than 1 ns.

It is probable that, at once after the squeezing of the scanning plasma layer into the collapse zone, an intense inverse shock wave is formed on the surface of this zone. If a plasma is present in the concentric slit between the collapse zone and the external layers of the target, this wave will propagate along it from the center to the periphery of the target. It seems to be quite probable that this inverse shock wave can stimulate the running of repeated nuclear reactions. It is worth noting that approximately the same (but less efficient) mechanism of the formation of an inverse shock wave was considered in many scenarios of the gravitational collapse of a massive star after the squeezing of the central iron-nickel core of the star. It is natural that this analogy is clearly formal, because the action of a gravitational field is absent in the above-analyzed process of collapse of a shell propagating to the center of a target. Such a difference is very important. In fact, the action of gravitation hampers the propagation of the inverse shock wave in the case of a gravitational collapse, whereas such a deceleration is lacking in our system.

It seems almost obvious that superheavy nuclei formed in the collapse zone for a sufficiently short period of the inertial confinement of the collapse zone from a free dispersion have no time to evolve up to a steady state at the expense of the realization of a long chain of successive nuclear interactions with nuclei from the composition of their nearest environment. In particular, according to the Migdal model, such a state of a nucleus can correspond to the mass number $A \approx 200\,000$. This circumstance leads us to a very important conclusion that the increase in the mass of such nuclei can also continue after their ejection from the collapse zone and the penetration into the volume of structural elements of the experimental setup surrounding the target. The analysis of the processes of nuclear transformations in these structural elements (named conditionally as accumulating screens) allows us to identify, in a clear manner, superheavy nuclei and to determine their certain parameters. In the following section, we will carry on the quantitative

simulation of the processes related to these transformations and observed in the executed experiments.

Further, we intend to perform a comprehensive study of the possibility and the features of the running of collective nuclear reactions in the zone of the scanning layer and in the zone of action of the spherical collapse.

The above-discussed scenario, despite the necessity of further corrections, allows us to explain the majority of the results of a lot of experiments on both the shock compression of a substance and the synthesis of heavy and superheavy nuclei and isotopes with anomalous ratios.

Analysis of the Spatial Distribution of Synthesized Isotopes in Accumulating Screens Remote from the Collapse Zone and the Problem of Evolution of Superheavy Nuclei. The above-considered scenario of the stage-by-stage “scanning” synthesis allows us to explain the possibility for the attainment of a superdense state of a target without any initial external threshold compression able to ensure the conditions for a collapse of the target substance. A natural consequence of the formation and evolution of such a collapse is the synthesis of a great amount of various isotopes and the formation of superheavy nuclei. The shift of the binding energy maximum for such nuclei to the region corresponding to very large mass numbers of nuclei allows us to conclude that the increase in the mass of such nuclei is an energy-gained process. Therefore, newly created nuclei are stable relative to the process of spontaneous fission. Since the electron shells around such nuclei should cardinaly differ from the standard “Bohr” atomic shells, the direct observation and identification of such supermassive quasiatoms by using the analysis of the spectrum of a characteristic emission are a complicated problem.

Below, we present the analysis of the experiments which reveal the anomalous spatial distribution of many elements and isotopes in the volume of an accumulating screen remote from the target. This analysis presents though an indirect, but sufficiently grounded confirmation of the existence and evolution of such superheavy nuclei. In addition, the simulation and calculation of the motion and interaction of such nuclei with atoms of the substance of an accumulating screen allow us to determine their parameters.

The analysis of the layer-by-layer distribution of chemical elements in the volume of a chemically pure accumulating screen reveals the following main regularities of the distribution of the concentrations of these elements.

The distribution of elements and isotopes was studied with the help of an ion microprobe IMS 4f (CAMECA, France) upon the ion etching of the surface of an accumulating screen earlier subjected to the pulse action from the side of the collapse zone. This investigation was carried out in

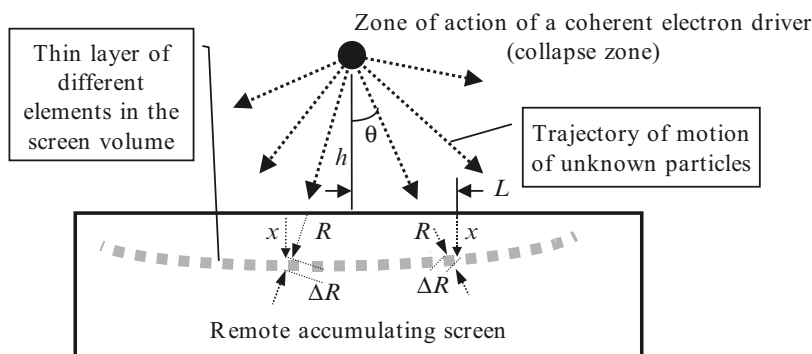


Fig. 11.24. Mutual position of the collapse region, remote accumulated screen, and layer of chemical elements and isotopes with anomalous spatial distribution in the screen volume; L is the distance from the screen center defined by the position of the collapse zone projection on the screen surface to a specific point of observation on the screen surface, h is the height of the collapse zone above the surface of the accumulating screen.

a sufficiently large time interval after the very process of formation and destruction of a collapse (as a rule, in the scope of 1–2 weeks).

Having performed a comprehensive study of the structure and composition of accumulating screens manufactured of a chemically pure element (with a purity of 99.99 to 99.96 %) and positioned near the target (see Fig. 11.24), we discovered, in many cases, the unique spatial distribution of various chemical elements and isotopes in the interval of atomic masses $1 \leq A \leq 240$.

The analysis with the help of an ion microprobe allows us to discover the presence of several layers with enhanced concentrations of various elements and isotopes which were absent in the screen and target materials (or, at least, were present in the target and the screen as admixtures at concentrations which were less by many orders than those in the volume of the layers under study). Some results related to various experiments are given in Figs. 11.25 and 11.26.

On the surface of the accumulating screen, we observed the maximum concentrations of all extrinsic (i.e., differing from the materials of the target and the accumulating screen) chemical elements. The thickness of this first layer was about 100 to 200 ρA . It contained about 3×10^{18} atoms. These atoms were distributed over the entire surface of the accumulating screen (i.e., in the scope of all solid angles relative to the collapse zone), their concentrations decreased with increase in the angle θ with respect to the direction corresponding to the relevant least distance from the collapse region to the screen (i.e., upon the departure from the collapse zone). With

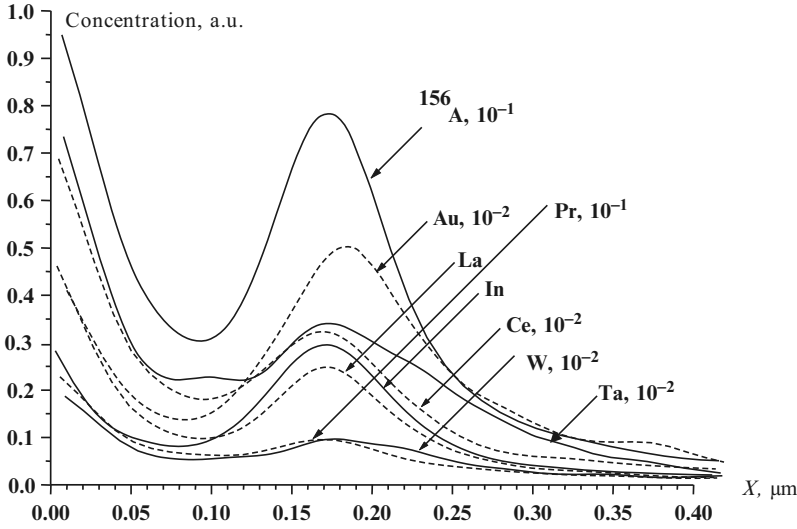


Fig. 11.25. Profile of the distribution of the concentrations of various chemical elements over depth in the region adjacent to the center of an accumulating screen.

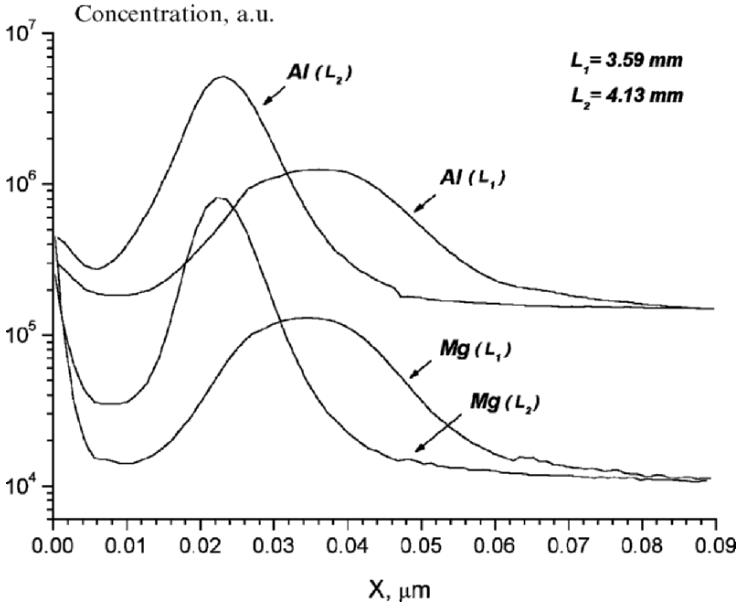


Fig. 11.26. Profile of the distribution of the concentration of Al and Mg for two different distances ($L_1 > L_2$) from the center of an accumulating screen versus the distance from the surface inwards the volume of an accumulating screen.

increase in the distance from the surface inwards the screen, we observed a fast decrease in the concentrations of these elements (see Fig. 11.25).

At a greater depth, we found the second layer, in the scope of which we registered the spatially coincident maxima of the concentrations of various chemical elements and isotopes. Depending on the conditions of the execution of an experiment and on the types of a target and an accumulating screen, the depth of these maxima at the points of observation (at a distance $L = 2$ to 5 mm from the conditional axis defining the projection of the collapse zone on the plane of an accumulating screen) is 0.02 to 0.2 μm under the surface of a screen. The evaluations show that this layer contains about 10^{18} atoms.

In the process of ion etching, we found that the position x_i of the distribution maxima of the concentrations of chemical elements varied by the law of inverse proportionality as a function of the distance from the screen axis: with increase in the distance L_i from the axis, the depth of the position of maxima x_i decreases. This result is seen in Fig. 11.26, where we present the distribution of concentrations of Al and Mg versus the distance from the screen center.

It is seen from Fig. 11.26 that, with increase in the distance from the screen center from $L = 3.59$ to $L = 4.13$ mm, the depth of the maximum of concentrations decreases from $x = 0.036 \mu\text{m}$ to $x = 0.024 \mu\text{m}$. In this case, the positions of the relevant maxima of the concentrations of various chemical elements completely coincide. It is also obvious that the concentrations of these elements on the surface are less by several times than those in the scope of a given layer. The effect of the increase in the registered half-width of the distribution of chemical elements near the distribution maximum position with increase in the depth of this maximum does not mean that the real half-width of the distribution grows as well.

This effect is related to the specificity of the process of ion etching of the surface. The point is in that the process of ion etching is not strictly deterministic, but is subjected to statistical fluctuations. In other words, upon such an etching, the surface of the etched layer is not ideally flat at any moment of the process of etching, but possesses a fluctuating profile and a nonzero variance and is characterized by deviations of both signs from the midplane. For this reason, even if the extremely thin layer of extraneous elements is present in the material under study, it will be registered as a layer with enhanced variance. The absolute value of the variance increases with the depth of the etched layer, which leads to the illusory growth of the distribution half-width for registered elements in that part of the layer which is located at a great depth. In fact, the real variance of the distribution of chemical elements and isotopes in the scope of a layer under study is very small.

On some screens, we discovered also the third concentration layer which is located at the depth of about 5 to 7 μm .

It is obvious that the surface distribution of chemical elements and isotopes in the first (surface) layer admits a natural explanation and is related to their direct arrival from the collapse zone at the screen surface under the destruction of the target. A decrease in the concentrations of various elements and isotopes upon the movement from the surface to the depth of this layer is well described by the ordinary law of diffusion.

Basically different are the origin and the interpretation of the registered maximum of the concentrations of various elements and isotopes in the scope of the second layer located in the volume of the accumulating screen.

The sequential analysis of regularities of a change in the concentrations of various elements and isotopes in the scope of the mentioned maximum leads to a number of paradoxical conclusions which cannot be explained without the attraction of a model of synthesis and evolution of superheavy nuclei. Below, we consider these conclusions.

Conclusion 1. It is easy to verify that the maxima in the volume of accumulating screens are not related to the process of diffusion of particles arrived from the outside. Indeed, if this were resulted from diffusion, the structure of the spatial distribution of a diffusing substance would depend only on the distance from the surface and would be independent of the point of the arrival of a diffusing particle at this surface (independent of the distance L from the axis). In this case, the distances x_i from the surface to all points of the second concentration maximum would be identical for all L_i . Moreover, diffusion is characterized by the monotonically decreasing distribution of the concentrations without maxima upon the departure from the surface. But, in the case under consideration, we registered the clearly pronounced maxima of concentrations.

In addition, we note that the diffusion coefficient depends very strongly on the characteristics of a diffusing particle (on its mass and charge) (this is, possibly, the principal point). The diffusion coefficient for heavy and light atoms can differ at least by tens of times. In this case, each type of atoms or ions would have the own law of distribution over the depth. However, we observed experimentally the same distribution of the concentrations of quite different elements (from hydrogen to lead)!

Conclusion 2. It is easy to verify the complete failure of the model, according to which the registered maxima of the concentrations are a consequence of the deceleration and the stop of the rapidly moving atoms (ions) of chemical elements and isotopes which came in the volume of an accumulating screen as rapidly moving particles.

It is well known from the physics of the passage of particles through a substance that the law of distribution of the path lengths with small variance is typical of fast heavy ions (i.e., one observes a clearly pronounced maximum of their concentrations).

Heavy ions, by supplying energy to atomic electrons of the target (whose mass is less than the mass of heavy ions by many orders), lose energy by very small portions in the form of ionization losses and stop after the full energy transfer.

The transfers of energy and momentum from an ion with mass M and kinetic energy E to an electron with mass m_e during one collision are small and equal

$$\Delta E/E \approx \Delta p/p \approx (m_e/M) < (10^{-4} \text{ to } 10^{-6}). \quad (11.149)$$

For this reason, heavy ions decelerate gradually by preserving the almost rectilinear trajectories. In this case, the variance of path lengths is small and does not exceed 2% to 4%. Thus, the position of a layer reflects the real distribution of path lengths of heavy particles along the initial direction of their motion.

It is easy to see from the analysis of the results of experiments that the different depths of this layer measured from the screen surface correspond to that the region with a concentration maximum in the scope of the layer is, in essence, the surface, every point of which is located at the same distance from the screen surface, if we measure this distance along the line connecting this point with the region of the collapse zone. In other words, let us assume that the collapse zone emits the identical particles with the same energy. Then, having come into the screen volume and having undergone the successive identical deceleration, they will stop at points which define, on the whole, the surface of the layer with high concentration of these particles. This surface is schematically drawn by the dashed line in Fig. 11.24.

This result follows directly from Figs. 11.24, 11.26, and the relation

$$x/\cos\theta = R = \text{const.} \quad (11.150)$$

This relation connects the local depth of the layer with the direction to the collapse zone and is well satisfied in the whole interval of values of x and θ .

Thus, the following situation is realized: the whole distribution of maxima can be explained by assuming that it is caused by the deceleration of particles with the same mean path length R .

This result seems to be quite paradoxical.

It is easy to verify that there exists no situation where fast ions of different types have the same path length.

For this purpose, it is necessary to take into account that the deceleration of fast nonrelativistic ions with mass M and charge Z in a substance with density of atoms n and with charge z is described by the Bethe-Bloch formula

$$dE/dr = -(2\pi n M Z^2 z e^4 / m_e E) \ln(4m_e E / M J). \quad (11.151)$$

Here, E is the energy of a particle, and J is the mean ionization potential of atoms in the target.

It is seen that, at the same energy E of a particle, the deceleration rate (the energy loss per path unit, dE/dr) depends on the mass M and the charge of a particle. These values define also the length of a full deceleration path

$$R = \int_E^0 dE / (dE/dx). \quad (11.152)$$

For example, in order that different ions pass the same distance $R = 0.3 \mu\text{m}$ in a copper target, it is necessary that each ion have a strictly fixed energy, whose value lies in the interval from 60 keV to 100 MeV and more (see Table 11.3).

Table 11.3. Kinetic energy of ions of various chemical elements for the path length $R = 0.3 \mu\text{m}$ in the volume of a copper screen.

Ion	E_K
$\text{H}^+ = \text{p}$	60 keV
Li^+	160 keV
Li^{++}	0.56 MeV
Be^+	0.23 MeV
Be^{++}	0.9 MeV
C^+	0.32 MeV
C^{++}	1.7 MeV
N^+	0.67 MeV
Ni^{++}	1.9 MeV
Pb^+	≈ 80 MeV
Pb^{++}	≈ 260 MeV

This requirement cannot be satisfied under any, quite fantastic assumptions.

If we assume that the particles emitted from the collapse zone to the space between the collapse zone and the screen surface move independently (this occurs, for example, at a very low density of the dispersed plasma), then their kinetic energy will be approximately the same and be defined by the temperature of the common source (the temperature of the collapse zone at the moment of its destruction). Their deceleration lengths will differ by several orders. This assumption contradicts the requirement that the energies should be different and correspond to the data presented in Table 11.3.

If we assume that the particles move in the volume of the dense plasma and arrive at the screen surface namely in such a way, then their mean linear velocity should be approximately the same. The direct measurements of the process of dispersion of a plasma flare outgoing from the collapse zone showed that this velocity is at most 2×10^7 cm/s. This velocity corresponds to the interval of kinetic energies ($E_K \approx 200$ eV for protons and $E_K \approx 40$ keV for uranium ions) which are less than the values corresponding to the path lengths of the particles presented in Table 11.3 by many orders. The analysis of the composition of the first (surface) layer confirms that the energy of ions corresponds approximately to the above-mentioned interval.

Moreover, since the efficiency of the deceleration of ions depends strongly on their charge, the ions of one element with different multiplicities of ionization will be decelerated in different ways and should be localized at different distances from the surface.

Conclusion 3. The derived distribution of the concentrations of various chemical elements is strange and has no obvious explanation. In Fig. 11.24, we show schematically that the surface of an accumulating screen is flat. In the initial state (prior to the realization of the action of an electron driver on the target), it was really flat. However, after such an action, it took the form of a concave surface (like a crater) with a maximum depth of $14 \mu\text{m}$ near the center. This depth exceeds considerably the measured distance from the layer surface to the anomalous concentration layer. The reason for such final form of the surface is obvious and is related with the removal (evaporation) of a part of the screen surface under the action of the X-ray emission from the collapse zone. The second reason for the removal of a part of the surface is the ion etching with a hot plasma created after the destruction of the collapse zone and the ejection of the hot plasma. These peculiarities are discussed in the section devoted to the description of the results of experiments.

This allows us to draw a number of important conclusions.

If the registered distribution of the concentrations of various elements and isotopes in the screen volume were related to the deceleration of various fast ions arrived from the outside, then the layer of ions in the central part of the accumulating screen would be closer to the surface, than that on the periphery. This would be related to that a part of the surface in this region would turn out to be etched (up to the depth considerably exceeding the path length of particles!). Moreover, fast ions would stop in the volume with regard to the fact that a part of them enter the target earlier, i.e., when the initial flat surface of the target was not else damaged (or was slightly damaged).

The experimental data indicate that all ions of different elements are located at the depth which corresponds strictly to the identical path length and is counted off the final concave profile of the surface. In this case, the maximum distance from the surface to the layer corresponds to the central region of the screen. This can occur only in the case where, on the one hand, the action of the emission or particles of the plasma changes the surface by forming a crater at its central part, but particles of the plasma do not penetrate into the screen due to their relatively small energy. Then, particles of the other type arrive at the screen, penetrate into it after the modification of this surface by particles of the plasma, and cause no changes in the surface.

In this situation, one more paradox appears. Since the distance from the collapse zone to the surface of an accumulating screen did not exceed several millimeters, these unknown particles would have a very small velocity to realize this scenario. Only in such a way, they can arrive at the screen surface already after the full completion of the process of removal of the mass from the screen surface.

It is obvious that if these unknown particles would be ordinary ions, they would have a small energy at a small velocity. But, possessing a small energy and a small velocity, these particles cannot penetrate into the target!

It is impossible to solve this paradox by remaining in the frame of the model, according to which different particles penetrate into the target.

Conclusion 4. One more contradiction has arisen after the execution of a more detailed study of the peculiarities of the distribution of various elements and isotopes on the surface of the discovered layer with enhanced concentrations in the volume of the accumulating screen. The direct investigation with the help of an ion microprobe IMS 4f showed that various elements and isotopes on the surface of the layer with enhanced concentrations are gathered in the scope of small and overlapping clusters, rather than are randomly distributed.

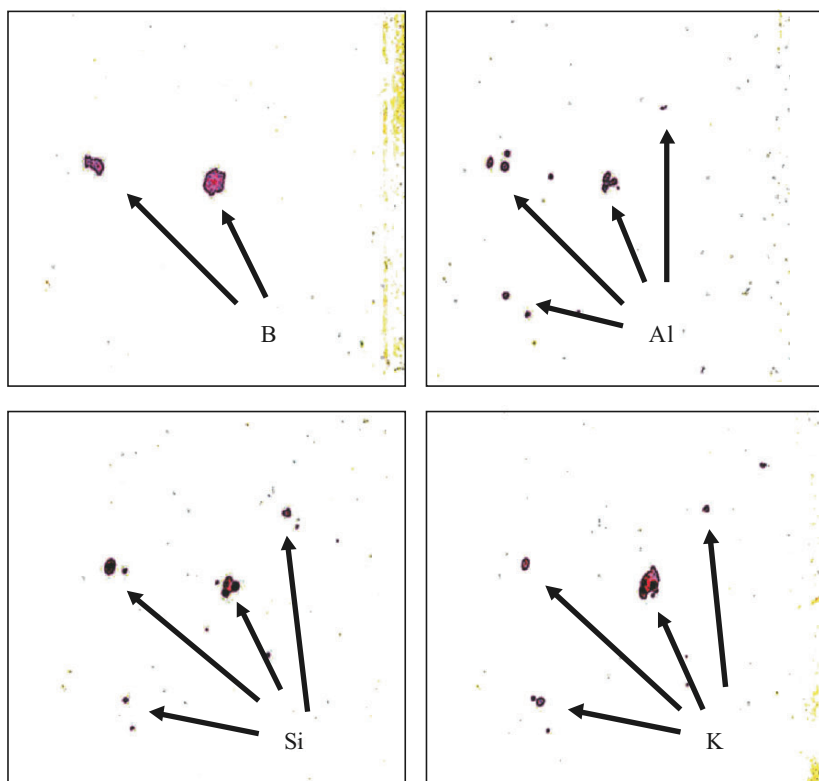


Fig. 11.27. Distribution of the clusters of chemical elements B, Al, Si, and K in the same very small imagining area on the surface of a thin layer.

In Fig. 11.27, we present the distribution of different chemical elements (B, Al, Si, and K) on the surface of the studied (etched) part of the layer in the form of a rectangular with its sides of about $100\ \mu\text{m}$. It is seen that different chemical elements are localized on the layer surface in the coincident regions! It is impossible to explain this result by basing on the viewpoint, according to which different elements and isotopes arrived at the given place of the layer after their ordinary linear deceleration.

All the above-presented facts and the phenomenon of the constancy of a depth and a location of the place of registration of various elements in the alien chemically pure matrix can be explained only on the basis of the following hypothesis. It is necessary to assume that a path with length R characterizes a slow (but not diffusive) motion of the strictly definite kind of identical unknown particles which have, in this case, a fixed charge (it is zero!). In addition, these unknown particles must have a very small velocity of motion and be very heavy. Only in this case, all the particles will be

braked in the scope of one layer equidistant from the screen surface along the direction of their motion from the collapse zone. In this case, all the set of registered elements and isotopes is a consequence of their creation (synthesis) upon the interaction of these stopped unknown particles with atoms of the screen.

Only in this case, the entire totality of the derived results can be satisfactorily explained.

Now we justify the correctness of the above conclusion.

If these hypothetical particles were fast (i.e., their velocity were more than $v_0 = e^2/\hbar = 2.5 \times 10^8$ cm/s typical of the valent electrons of atoms in the target), they would be ionized during their movement and, for this reason, would change their charge state. But it was shown above that the charge state of the observed particles must be the same during their deceleration. If these particles are slow (i.e., if they move with velocity v which is considerably less than v_0), they will remain always (in any medium) neutral. The particles must be slow by virtue of the other above-presented argument: they arrive at the screen surface after both the completion of the process of evaporation induced by X-rays and the ion-plasma etching of the surface, which can occur only in the case of small velocities.

To penetrate into the screen at a relatively great depth, particles must have a sufficiently high energy. With regard to a small velocity of motion, the mass of such particles should be very large.

For such slow, very heavy, and neutral particles, the Bethe–Bloch law Eq. 11.151 is not true (it was deduced only for charged particles).

Below, basing on the experimental data, we consider the dynamics of the motion of such unknown particles.

Their motion in the medium is similar to that of atoms (or molecules) in a gas (see Fig. 11.28).

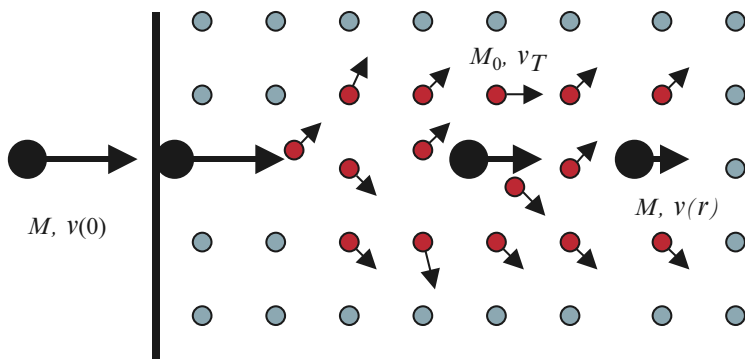


Fig. 11.28. The scheme of deceleration of unknown heavy particles in the volume of an accumulating screen.

The cross section of collisions of such particles with atoms of the target is defined by the radius of atoms R_0 . For a copper metallic screen, the scattering cross section is, with a sufficient accuracy,

$$\sigma = \pi R_0^2 \approx 10^{-16} \text{ cm}^2. \quad (11.153)$$

The mean path length of these particles between two subsequent collisions with atoms of the screen is very small and, at the concentration of atoms of the screen $n_s \approx 8 \times 10^{22} \text{ cm}^{-3}$ typical of solids, is equal to

$$l_1 = 1/n_s \sigma \approx 12 \rho A. \quad (11.154)$$

Hence, along the deceleration path length $R \approx 0.05$ to $0.3 \mu\text{m}$, there occurs a great number of collisions of these particles with atoms of the screen:

$$N = R/l_1 = R n_s \sigma \approx 40 \text{ to } 250. \quad (11.155)$$

In these collisions, a particle loses gradually its kinetic energy. Because every particle passes approximately the same total path and is not deflected sideways during collisions (all particles turn out finally at approximately the same distance from the input surface), it becomes obvious that particles lose a very small part of their energy in every collision, and their momenta are changed slightly. This means that the mass of particles M is considerably greater than the mass M_0 of atoms of the target.

At every point of the target, the motion of a particle is characterized by a definite velocity v which is quite gradually decreased along a trajectory of motion. Every collision of a particle with atoms of the target can be described in the system, where the particle is in the rest.

In this system, a target atom moving with velocity v (its momentum is $M_0 v$ in this system) collides with the immovable particle and then is elastically scattered. In every collision, the hypothetical particle varies the momentum of a target atom by the value $\delta p \approx 2M_0 v$.

The same momentum (but with the opposite direction) is received by the very particle for every act of collision.

The time interval between two successive collisions with the momentum transferred to the particle, δp , is

$$\delta t = l_1/v = 1/n_s \sigma v. \quad (11.156)$$

Thus, $\Delta N = \Delta t/\delta t$ collisions occur for the time $\Delta t \gg \delta t$, and the particle changes its momentum by the value

$$\Delta p = -\delta p(\Delta t/\delta t) = -(2M_0 v^2/l_1)\Delta t. \quad (11.157)$$

This relation yields the following formula for the mean braking force acting on the particle upon its movement across the target:

$$F \equiv \Delta p / \Delta t = -(2M_0 v^2 / l_1). \quad (11.158)$$

This force defines the law of motion and deceleration of the particle

$$M dv/dt = F = -(2M_0 v^2 / l_1). \quad (11.159)$$

A solution of this equation has the form

$$v(t) = v(0) / [1 + 2M_0 v(0)t / Ml_1], \quad (11.160)$$

where $v(0)$ is the initial velocity of the particle on the target surface.

It is obvious that this law of motion of the particle will be true as long as the kinetic energy of the particle $Mv(t)^2/2$ exceeds the thermal energy of atoms in the target $M_0 v_T^2/2 = (3/2)T_0$. With decrease in $Mv(t)^2/2$ up to $M_0 v_T^2/2$, the energy transfer from the particle to the target stops. The relation $Mv(\tau)^2/2 = M_0 v_T^2/2$ yields the duration of the process of deceleration

$$\tau = [(v(0)\sqrt{M}) / (v_T\sqrt{M_0}) - 1] Ml_1 / 2M_0 v(0). \quad (11.161)$$

The path length of the particle in the process of deceleration can be determined from the formula

$$\begin{aligned} R(\tau) &= \int_0^\tau v(t) dt = \int_0^\tau v(0) / [1 + 2M_0 v(0)t / Ml_1] dt \\ &= (Ml_1 / 2M_0) \ln \{1 + 2M_0 v(0)\tau / Ml_1\}. \end{aligned} \quad (11.162)$$

Taking τ from Eq. 11.159, we get

$$\begin{aligned} R(\tau) &= (Ml_1 / 2M_0) \ln \{ (v(0)\sqrt{M}) / (v_T\sqrt{M_0}) \} \\ &\equiv (Ml_1 / M_0) \ln \{ E(0) / T_0 \}, \end{aligned} \quad (11.163)$$

where $E(0)$ is the particle energy on the target surface.

Using the last formula, we can estimate the mass of a hypothetical heavy particle in terms of the main characteristics of the experiment [its path length $R(\tau)$ in the volume of the screen up to its stop, the density n_s , mass of atoms M_0 , temperature T_0 of the screen, and initial energy of the particle, $E(0)$] as

$$M = [4R(\tau)M_0/l_1] / \ln \{ E(0)/T_0 \} = 4R(\tau)M_0 n_s \sigma / \ln \{ E(0)/T_0 \}. \quad (11.164)$$

Below we make some numerical estimations.

At the time of the deceleration of a slow particle, the surface of the accumulating screen was sufficiently hot (after the removal of atoms from the surface). For the subsequent evaluation, we take the temperature of the screen $T_0 = 600$ K.

To calculate M , it is necessary to set a specific value of the particle energy $E(0)$ on the input surface of the screen. At first, we determine the value of M corresponding to the lower limit of the mass of the unknown particle. In this case, $E(0)$ can be equated to its maximum value equal to the temperature $T = Mv(0)^2/3$ of the source of particles (to the temperature of ions of the plasma in the collapse zone in the anodic needle) defining the initial velocity of the particle $v(0)$ on the input surface of the solid screen.

Let us take into account that the mean free path length of a slow particle in copper at $R \approx 0.3 \mu\text{m}$ is $l_1 = 1/n_s\sigma \approx 12 \rho A$. Then, taking $T_0 = 600$ K = 0.05 eV and evaluating the ion temperature of the dispersed plasma after the collapse occurred in the anode as $T \approx 35$ keV, we get the particle mass $M \approx 750M_0$, which corresponds, for a copper target, to the atomic number of a hypothetical nuclear cluster (a superheavy nucleus) $A_c \approx 4800$.

The initial velocity of motion of these nuclei upon the ejection from the collapse zone and on the screen surface is relatively small and equals

$$v(0) = (3T/M)^{1/2} \approx 3.5 \times 10^6 \text{ cm/s} \quad (11.165)$$

which is less than the velocity of valent electrons by 100 times.

The deceleration time for nuclei in the volume of an accumulating screen is sufficiently great and equals

$$\tau = [v(0)/v_T - 1]Ml_1/2M_0v(0) \approx 0.1 \mu\text{s}. \quad (11.166)$$

All the derived estimations are related to the experimental data of one typical experiment, whose results are given in Fig. 11.25.

To evaluate the parameters of such nuclei, we use their initial energy corresponding to the temperature of the collapse zone $T \approx 35$ keV. This value was chosen in view of the fact that namely such a temperature characterizes the maximum of the registered X-ray emission spectrum.

It is easy to verify that if the ion temperature will be lower, then the mass of a superheavy nucleus will be larger. In particular, if we take that the energy of these nuclei corresponds to the initial ion temperature in the collapse zone, $T \approx 1$ keV, then the mass number of the nucleus under study is $A_c \approx 6500$. At $T \approx 100$ eV, $A_c \approx 8500$. It is seen that the value of the initial energy of the unknown particle influences quite slightly the result of the determination of its mass.

For the other experiments, all the main features of the distribution of synthesized elements coincide qualitatively, but the very positions of the concentration maxima almost coinciding for different elements can be somewhat different. However, in any case, these concentration maxima correspond to the motion of superheavy nuclei neutralized by degenerate electrons.

Analogous estimations performed for the other typical experiment, whose results are presented in Fig. 11.25, lead also to a great value of the mass of the unknown particle. For both experiments, the deceleration path length in the volume of a screen $R \approx 0.05 \mu\text{m}$. Using formulas Eq. 11.164, we get that the mass of the unknown particle $A_c \approx 800$ to 1400 if we use the whole evaluating interval for the initial particle energy (from 35 keV to 100 eV). In this case, its initial velocity was

$$v(0) \approx 1.5 \times 10^6 \text{ to } 2.3 \times 10^6 \text{ cm/s.} \quad (11.167)$$

Thus, the sequential analysis of the performed experiments allows us to conclude that, as a result of the explosion, very heavy slow neutralized nuclear clusters with masses $A_c = 800$ to 7000 are emitted from the explosion zone and move with a small initial velocity $v(0) \approx 1.5 \times 10^6$ to 3.5×10^6 cm/s. These clusters can be characterized as superheavy nuclei neutralized by the compressed gas of degenerate electrons. At the very instant of the escape from the hot plasma region of the collapse zone, these particles can be charged (be ionized). But, upon the arrival at the surface of an accumulating screen, they turn out to be neutral. It is possible that they have captured atomic electrons in the outer cooler part of the collapse plasma or in the near-surface plasma before the screen surface. Most likely, they were neutralized on the surface of a solid screen. The neutral state of nuclear clusters will be preserved upon the movement in any cool medium.

The velocity of these particles turns out to be less than that of valent electrons of atoms in the target, $v_0 = e^2/\hbar = 2.5 \times 10^8$ cm/s, by 100 times. For this reason, these particles do not excite or ionize atoms of the screen upon the passage of the volume of the accumulating screen and, in addition, do not change the own charge state.

These particles slowly lose their velocity by transmitting their energy to atoms of the target in the process of elastic collisions, excite the phonon oscillations in the target, and eventually heat it. Being decelerated to the energy equal to the thermal energy of the target, the particles stop, being in the state of thermodynamical equilibrium with surrounding atoms. Because a very small energy is transferred at every act of collisions of a heavy particle with atoms, the process of deceleration occurs up to a great depth at the expense of the very large number of successive collisions, which decreases

the probability of fluctuations of the deceleration and corresponds to a small variance.

The duration of the flight of these particles with the velocity $v(0) \approx 1.5 \times 10^6$ to 3×10^6 cm/s from the collapse zone located in the target to a place of the stop in the volume of the screen is 1 to 3 μ s. It is obvious that their lifetime exceeds this value in any case. Such nuclei with mass numbers $A > 800$ are inexplicably stable by “the nuclear standards”. Indeed, according to all the known commonly accepted nuclear models, any nucleus with $A > 300$ (and, all the more, with $A \gg 300$) must decay for an essentially smaller time.

Moreover, basing on the above-mentioned models, we may assume that the decay of these superheavy nuclei decreases their binding energy and therefore they are stable objects in vacuum. For the confirmation, we indicate that the nuclei with masses in the interval $A_c = 800$ to 4000 were repeatedly registered by the method of Rutherford backscattering in analogous accumulating screens in several months after that they were undergone to the action from the side of the collapse zone on the same setup.

If the mass of these nuclei does not reach the value corresponding to the energy minimum (the maximum of the binding energy), then they are able to absorb ordinary nuclei being in the medium and are gradually transformed into absolutely stable nuclei. According to the calculations made by Migdal almost 30 years ago, stable superheavy nuclei with pion condensate can possess the maximum binding energy in the region of $A \approx 200\,000$. Such a character of the evolution of these superheavy nuclei is defined by that the dependence of their binding energy on the atomic number A has the second maximum. This question was discussed above in detail. Respectively, the dependence of the total energy on A has two minima (the “ordinary” minimum at $A \approx 60$ and the “anomalous” second one at $A \approx 200\,000$).

If the maximum of stability of nuclei at $A \approx 200\,000$ is really exist, then the superheavy nuclei created in the collapse zone will efficiently absorb the surrounding nuclei of a target with $A_M \approx 1$ to 200 and will increase their atomic numbers (in the limit, up to $A \approx 200\,000$). Such an absorption is easily realized due to the very high transparency of the Coulomb barrier in nuclei surrounded by a superrelativistic electron gas.

The energy excess formed upon this absorption of the target nuclei will be released by means of various channels (including the random ejection of “ordinary” nuclei from the volume of a superheavy nucleus rapidly growing in size). This process of cooling of nuclei has the fluctuation nature and can run by the following scenario.

After the capture of an ordinary nucleus, a superheavy nucleus is heated and reminds a boiling fluid. To release the excessive energy

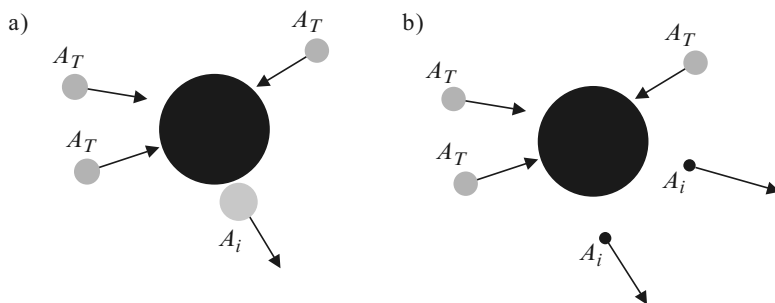


Fig. 11.29. Absorption of target nuclei and emission of one heavy nucleus (a) or several light nuclei (b).

(i.e., to cool a superheavy nucleus), some time is necessary. Therefore, if a superheavy nucleus absorbs simultaneously several nuclei of a target with $A_M \approx 1$ to 200, the excessive energy can be also released in the form of several light nuclei with $A_1 \approx A_M$ ejected from the superheavy nucleus or one heavy nucleus with $A_2 \approx 300$ to $500 > A_M$. This process was considered above in Sec. 11.1.4 and is schematically presented in Figs. 11.9 and 11.29.

It is obvious that the ejection (or evaporation) of newly formed nuclei is a competitive channel relative to other channels of the energy release from a growing superheavy nucleus. The process of evaporation of nuclei reminds the process of vaporization with the escape of molecules from a boiling fluid. Such a “reprocessing” of a substance surrounding every superheavy nucleus is completed upon the increase in its mass up to that corresponding to the state of full stability (up to $A \approx 200\,000$).

There are weighty reasons to believe that the anomalous distribution of chemical elements and isotopes registered in the experiments is the direct consequence of such an “activity” of growing superheavy nuclei. Isotopes formed in this nuclear reprocessing are registered immediately in the region of localization of particles. This result shows, besides other things, that the kinetic energy of synthesized and evaporated nuclei turns out to be small. The above-presented data on the braking lengths for various nuclei yield at once that the synthesized nuclei having a relatively small kinetic energy can go away from the region of their creation (i.e., from the zone of a full braking of initial superheavy nuclei) at a very small distance (less than $0.1\,\mu\text{m}$). This fact explains the identical widths of the regions of localization for different elements presented in Fig. 11.27. This result is also in a good agreement with the effect of the absence of radioactivity in the products of nuclear transformations.

The above-considered scenario confirms indirectly the presence of superheavy nuclei and explains the majority of regularities of the spatial localization of chemical elements and isotopes created, as a result of the experiments, in the volume of an accumulating screen positioned near the collapse zone.

NUCLEAR COMBUSTION AND COLLECTIVE NUCLEOSYNTHESIS

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and A. V. Paschenko**

12.1. Introduction: Collective Processes of Nucleosynthesis

The development of many directions of fundamental and applied nuclear physics is topical for many decades. As evidence of a great interest in the problems of nuclear physics, we mention the creation of such a big document as the Long Range Plan 2004 (see Ref. 208) by the Nuclear Physics European Collaboration Committee (NPECC) with the participation of about 300 leading experts in all branches of nuclear physics which defines the program of activity in this important scientific direction from the midterm perspective.

At present, the main problems of nuclear physics are recognized to be the clarification of the main structures of a nuclear substance and the principles governing the formation of stable structures in nucleons and nuclei and the expansion of our knowledge about superheavy nuclei and the limits of their stability. Experimental and theoretical investigations performed in the Electrodynamics Laboratory “Proton-21” and presented in this book are, in our opinion, a significant step in this direction.

At first look, even a simple enumeration of the effects observed in the Electrodynamics Laboratory “Proton-21” and related to the transmutation of nuclei (see Ref. 209) leads to the conclusion on the existence of the unremovable contradiction between them and the orthodox ideas of the contemporary nuclear physics. However, the careful analysis of the derived data and the fundamental ideas of the physics of condensed media allows one to destruct this first impression.

The first theoretical models leading to the conclusion about the rightfulness of the observed effects and the conception of self-organized nucleosynthesis have been developed in our laboratory and published in a number of works and preprints (see, e.g., Refs. 209–212).

In this chapter, we present the main elements of one of the models of the appearance of dynamical nuclear structures. This model can explain

many of the physical phenomena observed in the Electrodynamics Laboratory “Proton-21” in the field of the self-organized nucleosynthesis running as a result of the interaction of hard-current relativistic electron beams (REB) with a condensed medium.

The physical model of self-organized nucleosynthesis presented below in this chapter is based on the subsequent account of the nonlinearity of processes and the phenomena of collectivity, nonequilibrium, and coherence arising upon the interaction of the powerful concentrated sources of energy with targets.

We note at once that the collective processes in targets occur on such a level of the action on condensed media that the system cannot be considered to be equilibrium or ideal.

We list the main physical effects and phenomena occurring upon the intense action on the condensed medium under conditions of the experiment:

- Formation of the electron flux with magnetic self-insulation, focusing of a REB in a narrow ring region on the anode, and scanning of the target surface; enhancement of a symmetry of the squeezing of a target
- Nonequilibrium phase transitions under the action of concentrated flows of energy
- Increase in the density upon the development of various nonlinear processes in the target substance
- Development of instabilities in the dense electron-nucleus plasma of a target, in particular, the instability as for the appearance of two-fluid flows; the formation of self-consistent fields and quasiperiodic structures in a substance
- Change in the dispersive, polarization-related, and correlative properties of the substance and the nuclear system in nonequilibrium states and their influence on nuclear processes
- Enhancement of the role of memory in the system in connection with the peculiarities of correlations in it and the anomalous diffusion in the coordinate and phase spaces in states with a power law of decay of correlations
- Effect of the nonequilibrium properties of a medium and collective fields on the change in the properties of nuclear systems

The study of the listed physical mechanisms under conditions realized in the experiment allows us to formulate a qualitative pattern of the phenomena occurring upon the action of a powerful heavy-current REB on the target substance.

This pattern (scenario) consists in a sequence of phase transitions ensuring such states of the dense plasma of a target, in which the concentration of great energy densities on microscales becomes possible. The closer

the approaching of the density and spatial scales of a nonideal plasma in a target to the parameters of the nuclear substance, the greater the effect of these macroscopic parameters on processes in the nuclear substance. In other words, as a result of the action on the nuclear subsystem, its properties become dependent also on the macroscopic system parameters realized in experiments.

The scenario of the development of processes in experiments can be presented as follows:

- Interaction of the flows of electrons and a target. The appearance, as a result of the shock compression of a target, and evolution of a steepening nonlinear wave-shell, i.e., the appearance and evolution of the nonlinear density wave possessing a high symmetry and a small, as compared with the target size, characteristic spatial scale in radius.
- Interaction of the fluxes of ions and electrons in the dense plasma of a target upon the self-consistent consideration of the field equations and the hydrodynamic or kinetic equations for charged particles leads to the development of instabilities.
- The final stage of the development of arising instabilities is the transformation of the one-fluid mode of the evolution of a nonlinear wave in the two-fluid mode and the formation of plasma-field structures possessing small spatial scales and large self-consistent fields (approaching the nuclear ones). Thus, the wave-shell of density of the condensed neutral substance, which appeared at the beginning of the evolution, is transformed in a self-consistent nonlinear plasma-field object with complicated structure.
- The arising field changes the initial and boundary conditions in the equations of nuclear dynamics.
- Under these conditions, dynamical processes can lead to the appearance of nuclear objects with complicated internal structure (fractal nuclear clusters).
- These changes cause a modification of the binding energy in nuclei and a modification of nuclear structures. In other words, there appear collective nuclear structures, in which the Coulomb barriers of nuclei are collectivized, and the new stable states of great nuclear systems can arise.
- There arises the nuclear substance whose structure is essentially defined by electromagnetic fields on scales approaching the nuclear ones.

To correct the pattern of the processes running in a target, it is necessary to construct an adequate mathematical description of the physical

models of the interaction of a REB with a target. Since the evolution of the system is accompanied by the appearance of the densities of plasma and field rising with time, it is natural to consider them in the frame of the electrodynamic and statistical description of basic processes.

In Sec. 12.1, we present the analytic survey of the main known results which are necessary, in our opinion, for the construction of a theoretical description of self-organized nucleosynthesis and analyze the reasons of failures to realize the nucleosynthesis within traditional approaches. The main fragments of such a description are given in the following sections.

First, we present the general ideas of the means of a qualitative description of the macroscopic properties of media in equilibrium and non-equilibrium states and then will describe the main effects arising upon the evolution of the macroscopic parameters of a medium under the action of external concentrated flows of energy.

It is natural to consider, first of all, the known ideas of the nucleosynthesis, namely the processes of formation of nuclei in nuclear reactions. At present, it is clear that the processes of nucleosynthesis of the nuclei heavier than Li occur in stars on various stages of their evolution (see Refs. 213–216). Astrophysical processes in stars are responsible for the variety of elements and nuclei we observe in the universe. The same processes are the sources of a huge energy.

In general, the theory of nucleosynthesis developed up to now describes successfully the principal peculiarities of the abundance of nuclides in the solar system. According to the modern ideas, there exist four epochs of thermonuclear synthesis in primary stars:

1. The proton-proton cycle (the combustion of hydrogen) leading to the transformation of hydrogen in helium
2. Combustion of helium with the formation of carbon:

$$^4\text{He} + ^4\text{He} + ^4\text{He} \rightarrow ^{12}\text{C} + \gamma$$
3. Combustion of nuclei ^{12}C and ^{16}O with the formation of nuclei Mg, Ne, S, Si, and others
4. Combustion of nuclei Si with the formation of Ni and Fe

After the fourth thermonuclear epoch, the reactions of synthesis of nuclei terminate, because nuclei Ni and Fe have maximum specific binding energy. Therefore, the creation of nuclides heavier than iron should involve more complicated processes. According to the known theories, the reactions of synthesis of heavy nuclei are those of capture of neutrons. To realize the latter, very intense neutron fluxes are necessary, which are possible only upon the bursts of supernovas, the nuclei of galaxies, or similar objects.

Two following possibilities of the formation of nuclei heavier than iron are usually under consideration:

- Slow process of capture of neutrons by nuclei (*s*-process). The time interval between two subsequent captures of neutrons by a nucleus is greater than the lifetime of created isotopes with respect to β -decay. To realize the *s*-process, the neutron fluxes of $10^{15} \dots 10^{16}$ particles per second are required. The distribution of heavy nuclei observed in nature can arise as a result of the development of these processes during several thousands of years. In *s*-processes, no isotopes heavier than ^{209}Bi are created, because the lifetime of such isotopes is less than the capture time.
- In order to explain both the appearance of the nuclei heavier than ^{209}Bi existent in the nature and some anomalies of the abundance of elements, it is necessary to involve the *r*-process (the fast subsequent capture of many neutrons by a nucleus for the time less than the lifetime of the isotopes created upon every capture). The process is realized at very intense fluxes of $10^{27} \dots 10^{40}$ neutrons per second. These processes can occur for a very short time interval equal to tenths of one second or a few seconds.

However, the modern theory of nucleosynthesis cannot be considered as unique and completed, because many problems remain unsolved: the ratios of peaks of the recorded yields, anomalies of the contents of nuclides and elements in various astrophysical objects, existence of the “by-passed” nuclei, i.e., neutron-deficient nuclei which cannot be created in the *s*- and *r*-processes, etc.

The attempts to reconstruct the processes occurring in stars on the initial stages of their evolution such as the reaction of synthesis of lightest nuclei under laboratory conditions were started in the 1950s (see Ref. 217). As for the problem of practical realization of the processes of nucleosynthesis of heavy nuclei, it was not considered at all, though the theoretical (see Refs. 218–221) and experimental studies (see Refs. 222–223) concerning the nucleosynthesis of transuranium elements are intensively carried on at many scientific centers.

The unexpected breakdown in the region of self-organized nucleosynthesis has happened at the very end of the 20th century. We mean the experiments performed at the Electrodynamics Laboratory “Proton-21”. Details of the organization of these experiments and the description of their main results are presented below.

These experimental investigations gave a number of results exciting and hardly explainable from the viewpoint of the orthodox ideas. As most

unexpected results, we note the high stability of synthesized elements and the almost complete absence of radioactive isotopes in their composition (besides the very fact of the creation of macroscopic amounts of elements in the range from light to superheavy ones).

There are the weighty reasons to consider that the staff of the Electrodynamics Laboratory “Proton-21” has first realized the controllably initiated explosive physical process of nuclear regeneration of a macroscopic amount of substance and the creation (synthesis) of a wide spectrum of the stable isotopes of light, medium, heavy, and superheavy (transuranium) chemical elements. The whole totality of many experimental data testifies to that the process of nucleosynthesis is self-consistent, and the energy necessary for its initiation is less than that autonomously produced during a transformation of the substance by many orders.

The results derived in these experiments are so unexpected that they require a special analysis of the known fundamental ideas and nuclear technologies for their theoretical substantiation. We mention the earlier data related to the regeneration of nuclei at low energies, e.g., works (see Refs. 224–225). We also mention the works on sonoluminescence (see Ref. 226), sonosynthesis (see Ref. 227), cold synthesis (see Ref. 228), and muonic catalysis (see Ref. 229) close to the theme under study in view of the importance of collective effects on various levels of the structure of a substance.

In some experiments on the irradiation with electromagnetic nanosecond-pulses of the aqueous solutions of salts and melted metals, a transmutation of elements was observed (see Ref. 225). The experimental investigations of the problems concerning the transformation of elements are refined and are extended all the time. However, the number of theoretical studies of these processes is considerably lesser. All this defines the topicality of the theme of our investigation.

Practically from its development (beginning from the works of Bohr and Wheeler), theoretical nuclear physics uses the ideas of macroscopical physics (physics of continua and statistical physics). Since the middle of the 20th century, the main tools of theoretical investigations in this field are methods of the many-particle theory. At the same time, experimental nuclear physics uses only the methods of high-energy particles based on binary collisions (as distinct from the majority of natural astrophysical phenomena based on collective processes).

In traditional approaches, collective nuclear processes are considered only in the nuclear combustion wave propagating from the target center, where the ignition of a reaction is realized, to the periphery.

In the new approach to the consideration of nuclear transformations in substance developed in Ref. 209, many-particle interactions and, what is

the main point, self-consistent electromagnetic fields causing the formation of nuclear clusters (clusterization) in collective nuclear reactions are taken into account already on the stage of initiation of the ignition.

As mentioned above, the present investigation is an attempt to develop such theoretical model which would allow one to clarify the evolution of states of a target upon the realization of self-organized nucleosynthesis with the use of a new mean of the shock compression of a substance (the initiation of a self-developing collapse in its bounded volume) and would be, in our opinion, the base of the formation of a general approach to the development of new nuclear technologies.

Below, prior to the consideration of new models and solutions, we present a brief survey of the known results and formulate the main elements of the proposed approach.

12.1.1. Key Problems of Inertial Nuclear Synthesis

A peculiarity of the interactions in the system of nuclei is the fact that these interactions are mainly defined by two types of forces: the strong interaction between nucleons and the electromagnetic interaction between protons. In this case, the characteristic spatial scales of observed manifestations of these interactions are strongly different: the strong interaction (attractive force) is short-range with a characteristic length of 2 fm, and the Coulomb interaction between protons (repulsive force) is long-range (see Ref. 230).

Coulomb repulsive forces of protons in different nuclei being at large (on the nuclear scale) distances hamper the synthesis of new nuclei from the initial ones which occurs by means of the repacking of nucleons. To realize the synthesis, the Coulomb barrier should be overcome or be neutralized.

As usual, it is concluded that the best natural way to overcome a Coulomb barrier is to supply two nearest nuclei with the energy of the order of the height of a Coulomb barrier. This would ensure their approach up to a distance of the order of several fm, which would lead to their fusion.

In this case, the main problem, which has to be solved to realize controlled nuclear processes, is that of a self-consistent delivery of energy on nuclear scales in the minimum amount sufficient for the realization of nuclear transformations. The amount of the energy remaining in a compound nucleus after the overcoming of the Coulomb barrier should not strongly heat the nucleus and should be much less than its binding energy in order to ensure the stability of the compound nucleus.

Processes possessing the above-mentioned properties can be realized only in the collective self-consistent way and are the example of collective nuclear processes. Such collective nuclear transformations can be represented as a repacking of nucleons for the formation of new dynamical nuclear

structures in the relatively cool nuclear substance with a temperature lesser than the characteristic nuclear energies of the system.

The problem of energy concentration on nuclear scales is very complicated because the scales of phenomena observed by the man in the ordinary life and the scales of the nuclear world are hugely different. Characteristic sizes in our everyday activity are usually measured in centimeter and meter, whereas nuclear processes occur on the scales of the order of several fm. That is, these scales differ by 15 orders. This difference is similar to that between the scales of our everyday life and the size of the solar system. For this reason, to deliver the energy on nuclear scales, it is necessary to employ a suitable physical mechanism.

The ordinary approach to the solution of this problem involves binary collisions of accelerated nuclei and nuclei of a target of chaotic collisions between nuclei of a target heated to great temperatures in order to realize the thermonuclear synthesis (see Ref. 217, 231). The use of the technology of binary collisions is based on the idea of that individual processes are significant in nuclear transformations.

Such an approach does not solve, in fact, the posed problem completely, since the above-presented condition of consistency of the energy transferred to nuclei and the energy corresponding to the Coulomb barrier is not satisfied. Upon binary collisions, it is difficult to avoid the overexcitation of the compound nuclear system, i.e., its overheating.

The attempts to construct a thermonuclear installation with some force method of confinement of a nuclear fuel are also unsatisfactory. They deal with a macroscopic number of particles, but are based, in essence, on the same ideas of the defining and self-sufficient role of binary interactions.

In the scope of the traditional approaches to the realization of inertial synthesis, the confinement of a nuclear fuel happens at the expense of the own inertia of a substance. Most important becomes the problem of increasing the substance density upon its minimally necessary heating (it is desirable that the heating occur in the mode maximally close to the isentropic one) (see Ref. 217). The driver energy should be mainly expended not for the heating of the medium, but for the compression of a substance (see Refs. 231–232).

Upon the realization of the schemes of inertial synthesis, the compression of a substance occurs, as well known, at the expense of the complex processes of interaction of the energy flows with the surface. The main pressure is created by the evaporation (ablation) of the surface layer of a target at the expense of the transfer on the target surface.

The principal problem of thermonuclear synthesis is the realization of such a mode of the compression of a target, when only the central part

of a target is heated to the ignition temperature of reactions, whereas its peripheral part remains to be cold in this case. Such isoentropic compression is assumed to be realized with nonlinear convergent spherical waves which should be initiated as a result of ablation. After the initiation of the reactions of synthesis, one expects the propagation of a nuclear combustion wave from the center to the periphery through a cold substance, which would be accompanied by the energy release.

Thus, the key problem is to attain the superhigh compression of a cold substance. Consider this problem in more details.

For the approximate description of a target at a great pressure, we may use the model, in which the target is considered as a conducting fluid. In the one-fluid model, the hydrodynamic motions of a solid-state plasma do not differ from those of liquid metals.

The state of a medium is described with six quantities: three components of the velocity $\mathbf{u}(\mathbf{r}, t)$, the medium density $\rho(\mathbf{r}, t) = m n(\mathbf{r}, t)$, pressure $p(\mathbf{r}, t)$, and entropy $S(\mathbf{r}, t)$ [or temperature $T(\mathbf{r}, t)$]. In order to write the equations of the continuum hydrodynamics in the one-fluid mode, it is necessary to use the conservation laws for momentum, mass, and energy and the equations of state of a continuum. The conservation laws and the equation of state form the complete system of six equations for six listed variables.

We introduce the mass flow $\mathbf{j} = \rho \mathbf{u}$ and energy flow $\mathbf{j}_E$. Then the conservation laws can be written in the form of one vector and two scalar equations:

$$\frac{\partial \rho}{\partial t} + \text{div}(\mathbf{j}) \quad (\text{conservation of mass}), \quad (12.1)$$

$$\rho \left(\frac{\partial}{\partial t} + \mathbf{u} \nabla \right) \mathbf{u} - \frac{\rho}{m} \mathbf{F} = -\nabla P \quad (\text{conservation of momentum}), \quad (12.2)$$

$$\rho C_V \left(\frac{\partial}{\partial t} + \mathbf{u} \nabla \right) T + \text{div}(\mathbf{j}_E) = -\rho (\nabla \mathbf{u}) T \quad (\text{conservation of energy}). \quad (12.3)$$

In order to get a closed system, these equations should be supplemented by the thermodynamic relations defining the pressure P and energy flow. Usually, one uses the simplest equation of state and Fick's law:

$$P = \rho T / m, \quad \mathbf{j}_E = -\kappa \nabla T. \quad (12.4)$$

The one-fluid approximation is assumed to be valid by virtue of that the electron-ion equilibrium manages to be established in a cold compressed plasma.

The evaporation rate of a substance in the ablation zone is defined by the sound velocity in the hot region, and it is sufficiently high. The sound velocity in the cold region before the ablation zone is considerably lesser.

Therefore, the movement of the surface under the action of the ablation-related pressure is supersonic for a cold substance. A shock wave being a moving jump of the main parameters of a substance arises. Before the shock wave, the density, pressure, and velocity of the medium are equal to the initial values ρ_0 , p_0 , u_0 , respectively. Behind the wave front, the substance moves with the parameters ρ_1 , p_1 , u_1 .

The main properties of shock waves can be easily derived in the simplest ideal 1D case (a wave propagates along the x -coordinate) in a medium without viscosity and heat conduction. Equations 12.1–12.3 yield

$$\frac{\partial \rho}{\partial t} + \frac{\partial (\rho u)}{\partial x} = 0, \quad (12.5)$$

$$\frac{\partial (\rho u)}{\partial t} + \frac{\partial}{\partial x} (\rho + \rho u^2) = 0, \quad (12.6)$$

$$\frac{\partial}{\partial t} \left(\rho e + \frac{\rho u^2}{2} \right) + \frac{\partial}{\partial x} \left[\rho u \left(e + \frac{u^2}{2} + \frac{p}{\rho} \right) \right] = 0. \quad (12.7)$$

For the sake of convenience, we introduce the specific internal energy and specific enthalpy:

$$e = c_V T, \quad h = e + p/\rho. \quad (12.8)$$

Let a shock wave move with a constant velocity D . The propagation velocity of a shock wave is characterized by the ratio of its value to the sound velocity, $M_0 = D/c_s$, named by the Mach number. For a stationary motion, we may neglect the time derivatives in the hydrodynamic equations, integrate the remaining terms in the conservation equations (with space derivatives) along the wave motion direction, and get the algebraic relations for the jumps of the values before the discontinuity and after it (the relations of Rankine–Hugoniot):

$$\rho_0 u_0 = \rho_1 u_1, \quad (12.9)$$

$$p_0 + \rho_0 u_0^2 = p_1 + \rho_1 u_1^2, \quad (12.10)$$

$$h_0 + u_0^2/2 = h_1 + u_1^2/2. \quad (12.11)$$

From these relations, we get

$$u_0^2 = V_0^2 \left(\frac{p_1 - p_0}{V_0 - V_1} \right), \quad u_1^2 = V_1^2 \left(\frac{p_1 - p_0}{V_0 - V_1} \right), \quad (12.12)$$

where $V_0 = 1/\rho_0$ and $V_1 = 1/\rho_1$ are specific volumes. A change of the specific enthalpy can be written as

$$h_1 - h_0 = \frac{1}{2} (p_1 - p_0) (V_1 + V_0). \quad (12.13)$$

This relation and Eqs. 12.9–12.11 yield the functional relation between the pressures behind and before a shock wave, and the specific volumes. This relation is called the Hugoniot shock adiabat having different forms in dependence on the equations of state of a substance. In order to get closed formulas, we consider the simple case of an ideal gas with the polytropic equation of state

$$pV^\gamma = \text{const.}, \quad \gamma = c_p/c_V. \quad (12.14)$$

In this case, we have

$$e = c_V T = \frac{\rho V}{\gamma - 1}, \quad h = c_p T = \frac{\gamma p V}{\gamma - 1}. \quad (12.15)$$

Using these relations, it is easy to derive the relations for jumps of the wave parameters:

$$\frac{p_1}{p_0} = \frac{(\gamma + 1) \rho_1 - (\gamma - 1) \rho_0}{(\gamma + 1) \rho_0 - (\gamma - 1) \rho_1}, \quad (12.16)$$

$$\frac{\rho_1}{\rho_0} = \frac{(\gamma + 1) p_1 + (\gamma - 1) p_0}{(\gamma - 1) p_1 + (\gamma + 1) p_0}, \quad (12.17)$$

$$\frac{T_1}{T_0} = 1 + \frac{2\gamma}{(\gamma + 1)^2} \left(\frac{\gamma M_0^2 + 1}{M_0^2} \right) (M_0^2 - 1). \quad (12.18)$$

If we express the pressure in terms of the specific volume, Eq. 12.16 yields the analytic formula for the Hugoniot adiabat:

$$p_1 = p_0 \frac{(\gamma + 1) V_0 - (\gamma - 1) V_1}{(\gamma + 1) V_1 - (\gamma - 1) V_0}. \quad (12.19)$$

In Fig. 12.1, we draw the Hugoniot adiabat and the ideal gas adiabat under isentropic compression.

A shock adiabat is positioned strictly above the isentropic curve. Since the area under a curve in the coordinates (p, V) represents the energy, it is clear that one should expend much more energy to reach a fixed compression for the mode with shock waves than that for the isentropic mode.

In addition, let us express the medium parameters through the Mach number of a shock wave:

$$\begin{aligned} \frac{p_1}{p_0} &= 1 + \frac{2\gamma}{\gamma + 1} (M_0^2 - 1), \quad \frac{\rho_1}{\rho_0} = \frac{(\gamma + 1) M_0^2}{(\gamma - 1) M_0^2 + 2}, \\ \frac{T_1}{T_0} &= \frac{\gamma - 1}{\gamma + 1} \left(1 + \frac{2\gamma}{\gamma + 1} (M_0^2 - 1) \right). \end{aligned} \quad (12.20)$$

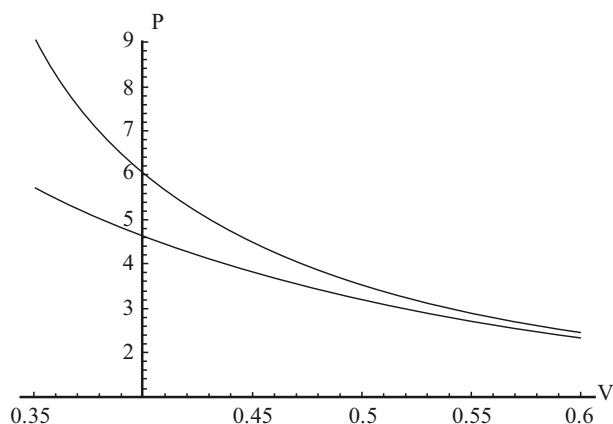


Fig. 12.1. Hugoniot shock adiabat and the adiabat for the isentropic compression of the ideal gas.

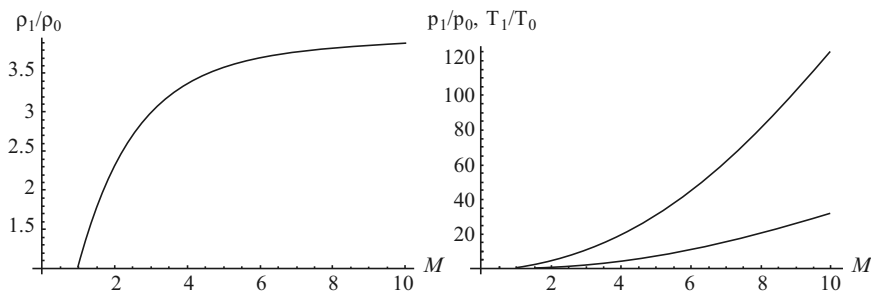


Fig. 12.2. Parameters of a shock-compressed substance versus the shock wave intensity.

As seen from Fig. 12.2, the pressure and temperature increase infinitely with the shock wave intensity, and the value of compression is bounded.

Due to the action of of an arbitrary large energy pulse on the surface of a target with the excitation of a strong shock wave, one may reach the compression of a target only by four times, if no efforts are undertaken for a specific organization of the experiment. As known (see Ref. 217, 231), the situation can be somewhat improved by means of the use of multilayered targets. In this case, at the expense of the specific sequence of the interface of the media, we can realize the compression by a sequence of weak shock waves imitating the isentropic process. However, the necessary driver energy turns out to be too great (see Ref. 217).

Up to now, we have considered plane shock waves. Let us see how the situation is changed, if the spherically symmetric compression of a substance is produced by spherical convergent waves.

The Euler equations of hydrodynamics for such 1D radial motion have the following simple form

$$\frac{\partial u}{\partial t} + u \frac{\partial u}{\partial r} + \frac{1}{\rho} \frac{\partial p}{\partial r} = 0, \quad (12.21)$$

$$\frac{\partial \rho}{\partial t} + \frac{\partial (\rho u)}{\partial r} + \frac{2\rho u}{r} = 0, \quad (12.22)$$

$$\frac{\partial S}{\partial t} + u \frac{\partial S}{\partial r} = 0, \quad (12.23)$$

$$p = p(\rho, S) \quad \text{or} \quad S = S(\rho, p). \quad (12.24)$$

The same equations can be written in the Lagrange form, if we consider the evolution with time t of the coordinates r of a particle of the continuum, being at the given initial point R at the given initial time moment t_0 . In the case of central symmetry, such equations have some advantages as compared to the equations for the Euler variables.

First, they describe the movement of every individual macroparticle and are convenient for the analysis of evolutionary processes. Second, namely in the case of spherical symmetry, the equations in the Lagrange form

$$u = \frac{\partial r}{\partial t}; \quad \frac{\partial u}{\partial t} + \frac{1}{\rho} \frac{\partial p}{\partial r} = 0, \quad \rho r^2 \frac{\partial r}{\partial R} = \rho_0 R^2, \quad p = \sigma(R) \rho^k, \quad (12.25)$$

where r is a running point and R is a value of r at $t = 0$, are simpler than the equations in the Euler form.

The question about the symmetric adiabatic compression of a finite mass of a substance by a self-organizing piston as a result of nonlinear hydrodynamic processes (the compression without appearance of shock waves) was strictly considered in (see Refs. 232, 233) on the basis of studies carried out in (see Ref. 234).

Up to now, we have discussed only 1D flows of a substance. This is related to the assumption of a high symmetry of the squeezing of a target. However, already in the case of a realization of the inertial synthesis in a cylindrical geometry, significant are 3D flows accompanied by the excitation of high-amplitude collective electromagnetic fields with several spatial components. For the one-fluid stage, the development of these processes is usually described by the equations of magnetic hydrodynamics (see Ref. 235) derived by Alfven for the first time. In this approximation, the usual equations of hydrodynamics are supplemented by the Maxwell equations for electromagnetic fields:

$$\operatorname{div} \mathbf{B} = 0, \quad (12.26)$$

$$\operatorname{rot} \mathbf{E} = -\frac{1}{c} \frac{\partial \mathbf{B}}{\partial t}, \quad (12.27)$$

$$\operatorname{rot} \mathbf{B} = \frac{4\pi}{c} \mathbf{j} + \frac{1}{c} \frac{\partial \mathbf{E}}{\partial t}. \quad (12.28)$$

The connection between two systems of equations in this approximation is realized through a magnetic field $\mathbf{B}$ and the current density $\mathbf{j}$. The right-hand side of the conservation law for momentum (Eq. 12.2) is supplemented by the Lorentz force $\mathbf{F}_B = (1/c) [\mathbf{j}, \mathbf{B}]$.

Magnetohydrodynamic processes in the dense plasma near a target and in the very target are essential for the dynamics of the whole process, because they induce a growth of the magnetic field in the system due to the development of various instabilities (see Ref. 236). We believe that the geometry of experiments in the Electrodynamics Laboratory “Proton-21” is significantly defined by the helicoid instability of a REB (see Refs. 236–237), in particular on the nonlinear stage of processes. Though the theory of these instabilities is developed, their role in the dynamics of electron-nucleus processes requires a separate investigation.

3D magnetohydrodynamic flows can possess the vortex and soliton-like structures. The role of these structures is very significant for the organization of a nuclear combustion wave in relation to their particular stability to considerable variations in their spatial scales. Stability of these structures is connected with the topological invariance of their states.

It is of interest that the exact solutions of such complicated system as the system of equations of magnetic hydrodynamics (even in the in two-fluid approximation) can be constructed on the base of Lagrange integrals and the freezing-in integrals without, in fact, solving this system (see Ref. 238).

We recall that Lagrange invariants I satisfy the equation

$$\frac{\partial I}{\partial t} + (\mathbf{v} \nabla) I = 0, \quad (12.29)$$

and a vector field $\mathbf{j}$ is called frozen in a fluid if it satisfies the equation

$$\frac{d\mathbf{j}}{dt} = (\mathbf{j} \nabla) \mathbf{v}. \quad (12.30)$$

In Ref. 238, the connection between these integrals is established, and a method of derivation of new freezing-in integrals from the system of Lagrange invariants I_k is developed:

$$\mathbf{j} = \frac{1}{\rho} \varepsilon_{\alpha\beta\gamma} \varepsilon_{ikl} I_i \frac{\partial I_k}{\partial x_\beta} \frac{\partial I_l}{\partial x_\gamma}. \quad (12.31)$$

In the general case of a 3D geometry in the two-fluid approximation, this method allows one to derive the vortex solutions of the system from the

analysis of the symmetry and conservation laws of the system without its direct solution.

Under powerful and nonstationary actions of concentrated energy flows on a solid, nonlinear effects have to reveal themselves and to become defining. We will show in the following chapter that the structural peculiarities of a diode used in the experiments carried out at the Electrodynamics Laboratory "Proton-21" lead to that the electron beam has two spatial components. One component propagates practically along the system axis, and the second one affects the lateral surface of a target.

The formation of such complicated structure of electron and ion fluxes near a target leads to the appearance of all components of the current and, hence, all components of the magnetic field. Moreover, a pulse of the magnetic field localized near the target surface is formed.

The pulse magnetic field concentrated on the lateral surface of the anode leads to the appearance of nonlinear waves of the magnetosound nature which propagate along a radius to the system center. Nonlinear magnetosound waves with high amplitudes were numerically studied in Ref. 239 and experimentally derived in plasma in Ref. 240. Work Ref. 241 presents a numerical and analytical generalization of results in Ref. 239 and shows that magnetosound waves with high amplitudes form a 2D quasi-soliton possessing a structure along both a radius and the system axis. The analytical solution of this problem turns out to be possible due to the introduction of Lagrange variables (r_0, τ) , in which the solution for a magnetic field can be written as

$$\frac{h(\tau, r_0)}{H_0} = T_h(\tau) s_h(r_0), \quad \frac{r(\tau, r_0) \omega_{pc}}{c} = T_r(\tau) s_r(r_0), \quad (12.32)$$

where the functions $T_h(\tau)$ and $T_r(\tau)$ are as follows:

$$T_h(\tau) = \frac{C_2}{C_3(\tau - \tau_0)^2 - 1}, \quad T_r(\tau) = \sqrt{C_1/T_h(\tau)}. \quad (12.33)$$

The function $s_h(r_0)$ in Eq. 12.33 is defined directly through the function $s_r(r_0)$:

$$s_h(r_0) = C_4 \frac{\frac{d}{dr_0} (s_r^2(r_0) ds_r(r_0)/dr_0)}{s_r(r_0) ds_r(r_0)/dr_0}. \quad (12.34)$$

The function $s_r(r_0)$ satisfies the equation

$$2 \left(\frac{ds_r(r_0)}{dr_0} \right)^2 - s_r(r_0) \frac{ds_r(r_0)}{dr_0} = C_5, \quad (12.35)$$

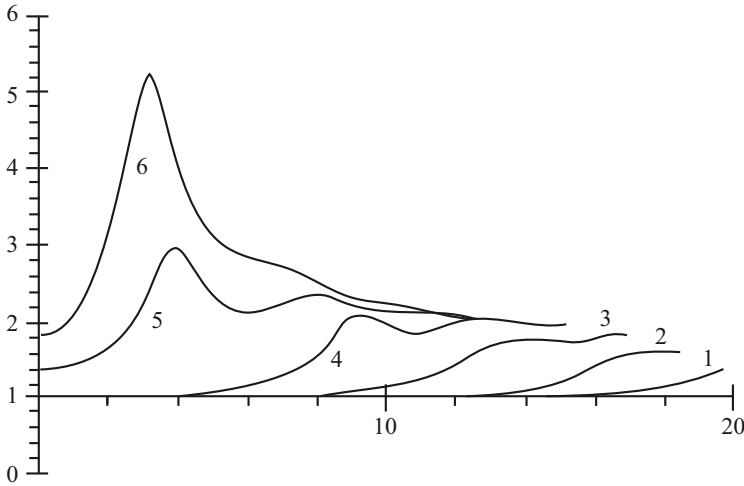


Fig. 12.3. Spatio-temporal evolution of a fast magnetosound wave. The dependence of the magnetic field of the wave on the radius for different times.

which is exactly integrable; we get

$$4C_5r_0 - C_6 = s_r(r_0) \left(-s_r(r_0) + \sqrt{s_r(r_0)^2 + 8C_5} \right) + 8C_5 \log \left(s_r(r_0) + \sqrt{s_r(r_0)^2 + 8C_5} \right). \quad (12.36)$$

As seen, we obtain the formula for a magnetic field through elementary functions (logarithm). With the use of algebraic transformations, this formula can be expressed in terms of cosh, as it happens in simple soliton-like solutions. In Fig. 12.3, we demonstrate the distribution of a magnetic field along the system radius for various times.

The region of localization of the magnetic field in this nonlinear wave becomes narrower, and the field amplitude increases upon the propagation along a radius. High values of the intensity of the pulse magnetic field together with electromagnetic oscillations excited in the medium can serve an efficient reason for the collective excitation of a nuclear system.

Nonlinear vortex formations can be created with the use of various drivers. In the next section, we show how nonlinear formations of the vortex type arise and evolve under the action of electron beams on the target surface.

The necessity to consider the nonlinear effects specific namely for solid-state theory was realized, of course, long ago. Indeed, thermal expansion of solids and their heat conduction cannot be explained if the anharmonicity of a lattice is neglected. Moreover, only weak nonlinearities were

taken into account (in the frame of perturbation theory), which led sometimes to the loss of important features conditioned by the nonlinearity of a problem.

It is of interest that the physics of strongly nonlinear phenomena enters the modern science in 1955 due to the Fermi–Pasta–Ulam paradox (see Ref. 242) related just to the study of pulse actions in the scope of simple models of the condensed state of a substance. By using a powerful computer, Fermi, Pasta, and Ulam analyzed the question on the excitation energy thermalization in nonlinear chains of atoms.

Numerical experiments revealed the effect of the absence of thermalization under the impact action on a nonlinear chain of atoms. In 10 years, this paradox was explained, and the term “soliton” now widely used in physics was introduced by Zabusky and Kruskal who connected the fact of the absence of thermalization in numerical experiments carried out by Fermi, Pasta, and Ulam with the presence of solitons in a nonlinear string. Really, active investigations of nonlinear effects in physics began just from that time (see Ref. 243). Many effects discovered in continua are also observed, in a specific manner, in solids: we mean, first of all, nonlinear effects in the electron plasma of solids.

The exact account of a strong anharmonicity of oscillations of a crystal lattice (see, e.g., see Ref. 244) can lead to new important consequences. For example, the current transfer in quasi-1D systems is realized with the help of nonlinear waves of charge density rather than by elementary excitations (i.e., by electrons). In polyacetylene, current carriers are solitons, and these solitons can be charged and have zero spin or have nonzero spin and be uncharged.

Modern successes of the inertial synthesis (see, e.g., Ref. 245) are connected with the conclusions of work (see Ref. 231), in which it was proposed to use the external action by an energy flow $q(t)$ incident on the surface in the sharpening mode, i.e., with the requirement of $q(t) \rightarrow \infty$ as $t \rightarrow t_f$. The typical view of a source for these modes is shown in Fig. 12.4.

Upon the irradiation of a target by a pulse electron beam, we may expect the formation of a centripetal wave of compression in it (see Ref. 246).

In the experiments on controlled thermonuclear fusion (CTF) with inertial confinement (see Ref. 247), the external layers of a multilayered spherical microtarget undergo ablation. Inside a microtarget, the density $\rho \propto 10^3 \text{ g/cm}^3$, pressure $P \sim 10^5 \text{ Mbar}$, and temperature $T \sim 10 \text{ keV}$ are recorded.

The sharpening modes leading to the compression of targets are connected with peculiarities of the strongly nonstationary and nonlinear processes, for which the phenomenon of metastable localization of various

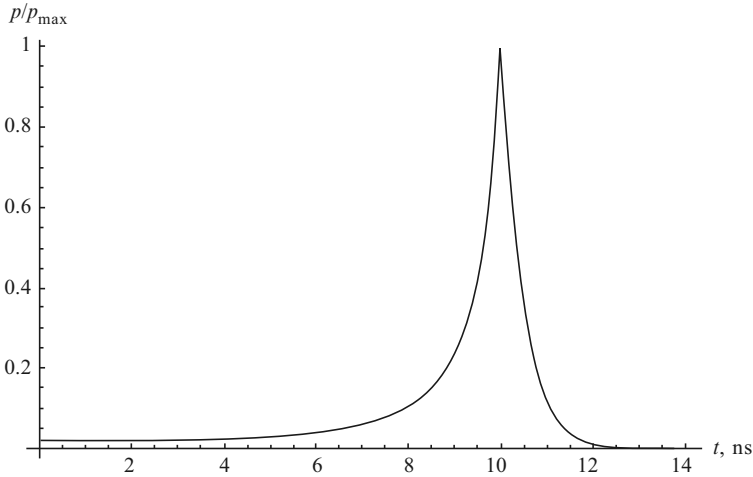


Fig. 12.4. Typical view of a pressure on the surface in the sharpening mode.

physical quantities in a nonlinear continuum is characteristic. In nonstationary systems, there occurs the limitation of the flows of main hydrodynamic quantities, which is described to the utmost adequately by the processes of their relaxation.

We demonstrate this by the example of the effects of energy (temperature) transfer. The transfer of temperature $T(\mathbf{r}, t)$ is defined by the energy conservation law in differential form,

$$\partial T(\mathbf{r}, t)/\partial t + \text{div}(\mathbf{j}_E) = 0, \quad (12.37)$$

where $\mathbf{j}_E$ is the energy flow.

As usual, it is assumed that the Fick law which connects the energy flow with temperature gradient is valid. The Fick law together with the energy conservation law lead to the parabolic equation of heat conduction and, hence, the infinite velocity of energy propagation (to be more exact, to the absence of the front upon the energy propagation in a substance). Long ago, Maxwell (see Ref. 250) called attention to the necessity to account the energy flow relaxation between collisions upon strongly nonstationary processes and wrote a more general relation for the flow with regard to its variation with a characteristic time τ_j :

$$\tau_j \partial \mathbf{j}_E / \partial t + \mathbf{j}_E = -\kappa \nabla T. \quad (12.38)$$

This relation together with the energy conservation law leads to the hyperbolic equation for temperature (the so-called hyperbolic equation of

heat conduction), which corresponds to a finite velocity of the energy propagation and the fronts of an energy density wave propagating from the surface source inward a target.

Such energy propagation can be accompanied by anomalous phenomena. For example, the relaxation of a heat flow admits the existence of an energy flow to the direction of increase in the temperature. In this case, thermodynamics is not violated since, under nonstationary conditions, the main inequality of the Prigogine's thermodynamics of irreversible processes (see Ref. 249) for the production of entropy σ_S ,

$$\sigma_S = \mathbf{j}_E \nabla \left(\frac{1}{T} \right) \geq 0, \quad (12.39)$$

should be replaced by the more general inequality

$$\sigma_S = (\mathbf{j}_E + \tau \partial \mathbf{j}_E / \partial t) \nabla (1/T) \geq 0. \quad (12.40)$$

Due to these thermodynamic anomalies, the processes with sharpening can cause the appearance and development of structures (regions with extremely high densities) in a medium.

If the certain conditions for an energy source on the surface are satisfied, the sharpening modes lead to the isentropic or almost isentropic evolution of the substance density.

In the case of inertial synthesis, of a great significance is the use of self-consistent nonlinear modes of compression of a quite cold substance. On the way of realization of the laser inertial synthesis, great successes in the compression of a substance with superpulses of the laser emission. We succeeded to increase the target density almost by three orders. However, despite the local successes, the problem was not solved, in fact. From our viewpoint, the main reason for failures is the incomplete use of the potentialities of collective processes upon the realization of nuclear transformations.

The action of the sources of concentrated energy flows on the target surface induces the appearance of states with very high density of a substance and sufficiently high temperature.

For the further analysis of the substance states with extreme parameters, it is necessary to have the idea of possible equations of state of the substance under critical conditions, of phase transitions, and of parameters of the system, at which these transitions occur. Below, we present a brief survey of the basic equations of state of a substance under conditions of the increase in its density and temperature.

12.1.2. Extreme States in Metals: Experimental Results and Limits of Theoretical Models

The correction of the limits of phase transitions in a substance for the region of extreme parameters and the analysis of dynamics of processes require a great volume of experimental data. The construction of the powerful sources of concentrated energy flows (lasers, the powerful sources of electron, ion, and neutron flows, shock and electromagnetic waves, etc.) has made the states of a superdense substance at earlier unattainable extremely high pressures and temperatures by the objects of laboratory investigations and applications in power engineering and technology. The available experimental data on the thermodynamical properties of dense substances are derived in dynamical experiments and generalized in reference books (see Ref. 251).

Some characteristic regions on the (n, T) plane are shown in Fig. 12.5.

The main problem of the theoretical description of extreme states of a substance consists in the presence of the strong interaction in a disordered medium, which excludes the application of perturbation theory to the

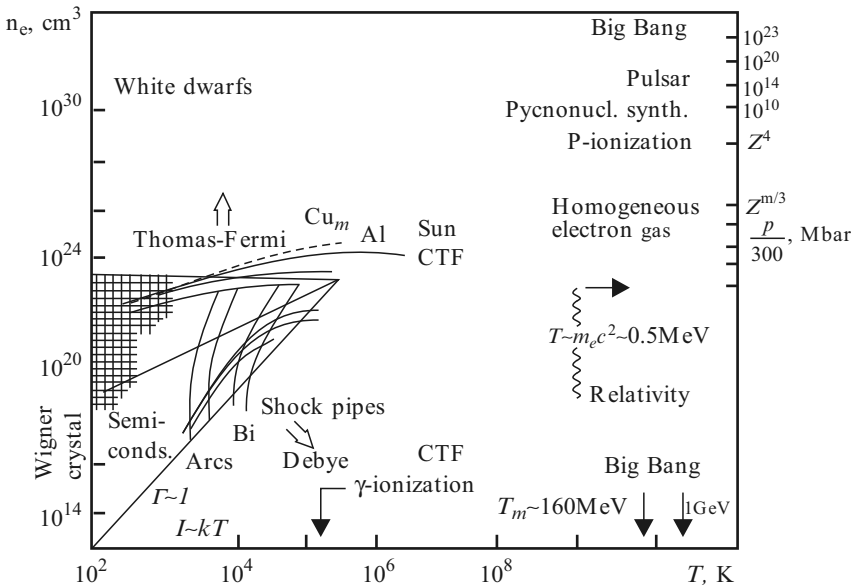


Fig. 12.5. Phase diagram of a substance. S_1 – adiabatic compression of saturated Cs vapors; $H'_1, H_2; H'_2, H_T$ – compression of saturated vapors of Cs and inert gases by incident and reflected shock waves; shock-wave compression of continuous – H_3 and porous – H_m metals; Al and Cu_m – shock adiabats of Al and porous copper; S_2 – adiabatic expansion of shock-compressed metals; Bi – isentropes of the expansion of Bi.

quantum-mechanical many-body problem. This circumstance makes it necessary to use physical models based on the simplifying ideas of the structure, energy spectrum, and character of the interparticle interaction.

In Fig. 12.5, we present the characteristic dimensionless parameters of a plasma, its technical applications, and the characteristic pressure (in the atomic system of units $P_a \approx e^2/a_B^4 \approx e^{10}m^4/\hbar^8 \approx 300$ Mbar) appearing in cosmic-physical and nuclear objects.

The relative value of the interparticle interaction in a Coulomb system is characterized by the dimensionless parameter $\Gamma = E_e/E_c$, which is defined by the ratio of the mean energy of the Coulomb interaction $E_e = Ze^2/r_e$, where r_e is the radius of screening, to the kinetic energy E_c .

Below the curve with $\Gamma \sim 1$, the interaction in the classical plasma is small and can be described within the model of chemical equilibrium (see Ref. 252), and the methods of perturbation theory (see Ref. 253) allow one to calculate the corrections related to the imperfection of the plasma. These calculations have asymptotic character, because they use the condition $\Gamma \rightarrow 0$. However, some appropriate ideas allow one to extrapolate these corrections onto the region up to $\Gamma \sim 1$.

The type of statistics of the electron component is defined by the degeneration parameter $n_e\lambda_e^3$, where $\lambda_e = (h^2/2\pi m_e k_B T)^{1/2}$ is the de Broglie thermal wavelength.

A characteristic scale of the kinetic energy is the Fermi energy $E_F = h^2 n_e^{2/3}/2m_e$, so that the compression of a plasma above the quantum imperfection limit $E_F \approx e^2 n_e^{1/3}$ leads to the simplification of its thermodynamical properties ($\Gamma \rightarrow 0$).

In Fig. 12.6, we give the phase diagram of Al with the indication of the approximate regions of validity of various theoretical models.

The model of ideal gas is the simplest approximation applicable in the region of small densities and not high temperatures lesser by approximately by the order than the first-ionization potential. In this case, the processes of thermal dissociation and ionization do not else occur in the system, and the interaction energy is small as compared to the kinetic energy of particles.

With increase in temperature, there occurs the dissociation of, first, the outer (at $T \propto 1 \dots 10$ eV) and then the inner ($T \propto 10Z^{4/3}$ eV) electron shells of atoms. The ionization equilibrium in such a system is described by the “chemical” model of plasma (the model of Saha (see Ref. 252)) based on the use of experimental and theoretical data of the excitation energies of atoms and ions and on their ionization potentials.

This model has a wide region of application bounded by the conditions of local thermodynamic equilibrium and stability ($T \leq m_e c^2 \approx$

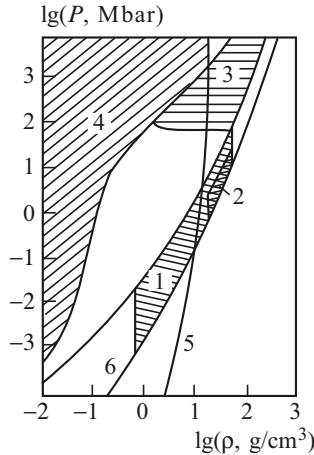


Fig. 12.6. Regions of validity of various approaches to the derivation of equations of state. 1 – semiphenomenological models; 2 – band model; 3 – model of Thomas–Fermi with corrections (TFC); 4 – model of ionization equilibrium (the Saha model); 5 – shock adiabat; 6 – isotherm ($T = 10^{-2}$ eV).

0.5 MeV) relative to the spontaneous creation of electron-positron pairs from the side of low densities and high temperatures and by the smallness of the interparticle interaction $\Gamma \ll 1$ from the side of high densities.

The efficiency of the quantum-mechanical methods of description of many-particle systems is defined in many respects by the choice of the zero approximation. The Hartree–Fock approximation is considered to be the best one-particle one (the self-consistent field approximation). However, the analysis of real systems in this approximation is associated with labor-consuming calculations and is not universal. One of the simple and simultaneously efficient models is the model of Thomas–Fermi (see Refs. 254–255) based on the quasiclassical approximation to the method of self-consistent field. This model is especially efficient under ultrahigh compressions, $P \gg P_a \approx 300$ Mbar.

The approximation of Thomas–Fermi was developed as a result of the solution of a particular problem of the structure of heavy atoms and at once became a tool theoretical physics used, in particular, in the description of a substance in extreme states. The relativistic generalization of this method, as will be seen below, is also convenient in the analysis of the processes in nuclei with regard to the processes of creation of electron-positron pairs in the field of a nucleus.

To clarify the role of collective and quantum effects under the Coulomb interaction, the model of one-component plasma against of the homogeneous neutralizing background of charges of the opposite sign is

widely used in the recent time (see Ref. 256). The results of a computer simulation of a one-component Boltzmann plasma under significant compressions demonstrate the appearance of a short-range order in the system, which was interpreted (see Ref. 257) as the formation of the amorphous phase at first and then the formation of a crystal lattice at $\Gamma \approx 178$. Phase anomalies in a degenerate electron gas were thoroughly considered in (see Ref. 258). In the course of the expansion of such a plasma, one observes the appearance of a charge-density wave and then that of a Wigner electron crystal (see Ref. 259).

At present, statistical models are in use for the derivation of the wide-range equations of state for solids due to the simplicity of their derivation and usage (see Refs. 253, 260).

Below, we consider the main theoretical positions necessary for the description of the processes of nucleosynthesis under extreme conditions.

12.1.3. Main Parameters, the Equation of States, and Phase Transitions of a Matter with Extreme Parameters

Objects, being the basis of physical systems we are interesting in, are charged particles (electrons and ions) obeying the Fermi statistics (fermions). The statistical properties of many-particle systems are of great importance for the understanding of physical processes. Below, we present the main relations of the statistics of fermions in equilibrium states.

Main Relations of the Statistics of a System of Identical Particles.

The basis for the description of a system of identical particles is their distribution function in the energy space. In an equilibrium state, the distribution over energy of fermions (for example, electrons) is defined by the Fermi-Dirac function (see Ref. 261)

$$f_{0e}(\varepsilon) = \left[\exp \left(\frac{\varepsilon - \mu_e}{T_e} \right) + 1 \right]^{-1}. \quad (12.41)$$

Here, T_e is the temperature of electrons, $\varepsilon = \varepsilon(p)$ is the energy of electrons, and the chemical potential of electrons μ_e is defined by the density of electrons through the normalization condition

$$\frac{2}{(2\pi\hbar)^3} \int_0^\infty 4\pi p^2 f_{0e}(\varepsilon(p)) dp = n_e. \quad (12.42)$$

A factor of 2 in the numerator is related to the number of states, $2s + 1$, with the same energy in a system of particles with spin s with regard to the fact that electrons have spin $s = 1/2$.

In the general case (with regard to relativistic effects), the dependence of the electron energy ε on its momentum p is given by the well-known formula

$$\varepsilon(p) = c \left(m_e^2 c^2 + p^2 \right)^{1/2}, \quad p = m_e c \left(\varepsilon^2 / m_e^2 c^4 - 1 \right). \quad (12.43)$$

Nonrelativistic and ultrarelativistic states are defined, respectively, by the conditions

$$p \ll m_e c, \quad \varepsilon(p) \approx m_e c^2 + p^2 / 2m_e, \quad p \gg m_e c, \quad \varepsilon(p) \approx cp. \quad (12.44)$$

By the Pauli principle, one state cannot be simultaneously occupied by more than one fermion. Therefore, while the temperature of electrons tends to zero, they try to occupy all energy levels from that with the lowest momentum up to the level with the maximum one p_F called by the Fermi momentum. Such a state is called degenerate. In a degenerate state, the distribution function is equal to 1 in the interval of momenta from zero up to p_F and is equal to zero for $p > p_F$. As a consequence, Eq. 12.42 yields

$$\frac{p_F^3}{3\pi^2 \hbar^3} = n_e, \quad p_F = \left(3\pi^2 \right)^{1/3} \hbar n_e^{1/3}, \quad \varepsilon_F = c \left(m_e^2 c^2 + p_F^2 \right)^{1/2} - m_e c^2, \quad (12.45)$$

and the Fermi energy ε_F is equal to the chemical potential:

$$\varepsilon_F = \mu_e. \quad (12.46)$$

In the limiting cases, the Fermi energy of electrons is defined in terms of their density by the simple relations

$$\begin{aligned} \varepsilon_F - m_e c^2 &\approx (1/2m_e) (3\pi^2)^{2/3} \hbar^2 n_e^{2/3} \quad \text{for } p \ll m_e c, \\ \varepsilon_F &\approx (3\pi^2)^{1/3} \hbar c n_e^{1/3} \quad \text{for } p \gg m_e c. \end{aligned} \quad (12.47)$$

The energy density of electrons in a degenerate state ρ_e is easily determined by values of the limiting Fermi momentum:

$$\begin{aligned} \rho_e &= \frac{E}{V} = \frac{2}{(2\pi\hbar)^3} \int_0^{p_F} c \left(m_e^2 c^2 + p^2 \right) 4\pi p^2 dp \\ &= 4K_e \left[x_F \left(1 + 2x_F^2 \right) \left(1 + x_F^2 \right)^{1/2} - \ln \left(x_F + \sqrt{1 + x_F^2} \right) \right], \\ K_e &= m_e^4 c^5 / 32\pi^2 \hbar^3, \end{aligned} \quad (12.48)$$

where $x_F = p_F / m_e c$. The pressure can be easily calculated from the relations

$$P = - \left(\frac{\partial (V\rho_e)}{\partial V} \right)_{N_e = Vn_e}. \quad (12.49)$$

In this case, we must take into account that the quantity x_F depends on the volume V . Thus, we get

$$P_e = \frac{4}{3} K_e \left[x_F (2x_F^2 - 3) (1 + x_F^2)^{1/2} + 3 \ln (x_F + \sqrt{1 + x_F^2}) \right]. \quad (12.50)$$

In these relations, it is assumed that the temperature is sufficiently small: $T \ll \varepsilon_F - m_e c^2$.

With increase in the density, i.e., with decrease in the mean distance between particles, atomic energy levels begin firstly to broaden out, which is conditioned by the tunnel effect. The volume of every atom is a potential well for the electron being in it. When the centers of atoms approach each other, the width of a barrier between adjacent wells decreases. For this reason, there appears the possibility for electrons to percolate from one atom to another one. The probability of the tunnel effect increases exponentially with decrease in both the barrier width and the difference of the barrier height and the particle energy.

The tunnel effect leads, obviously, to the additional broadening of energy levels. At a certain value of the density, an energy level broadens out so that it merges with the continuous spectrum. Then the electron in this state becomes free. It is obvious that optical electrons leave the atom in the first place, and electrons on the K -shell are the latest ones. The full ionization of atoms is realized in the case where the lowest energy level joins the continuous spectrum (see Ref. 262).

Let us estimate the density necessary for the full ionization of a substance. The phenomenon of percolation of electrons into adjacent atoms begins to efficiently act in the case where the mean distances l_* between atoms become comparable with the radius a_B/Z of the K -shell, where Z is the nucleus charge and $a_B = h^2/m_e e^2 = 0.529 \times 10^{-8}$ cm is the Bohr radius. That is, the density of atoms necessary for the full ionization is

$$n \approx \frac{1}{l_*^3} \approx \frac{Z^3}{a_B^3}. \quad (12.51)$$

If the density exceeds this value, the main share of the mass is in the state of full ionization. We will clarify now the state in which the electron gas is in this case: Is it relativistic or nonrelativistic, degenerate or not? The electron gas is relativistic if $p_F > m_e c$. With regard to the Fermi energy Eq. 12.47, we get the condition of relativity as

$$n_e \geq 0.1/\lambda_e^3 \approx 2 \times 10^{30} \text{ cm}^{-3}, \quad (12.52)$$

where n_e is the electron concentration and $\lambda_e = h/m_e c$ is the Compton wavelength of an electron.

The baryon substance in the liquid state (in ordinary nuclei) has density about $n_0 = 1.4 \times 10^{38}$ baryon/cm³. The mean kinetic energy of nucleons in such a liquid is of the order of 25 MeV, and the estimates of nuclear compressibility show that the sound velocity in the nuclear substance is close to 0.3 c.

The baryon substance can be considered on two levels, and each level has the own phase states.

The first level corresponds to the nuclear substance consisting of nucleons and possessing only the nucleon-related degrees of freedom. The second level corresponds to the baryon substance with quark-related degrees of freedom which reveal themselves to the full extent at high temperatures and densities, when a quark-gluon plasma is formed (see Ref. 263).

Between these levels, there exist a huge number of phase states.

Let us consider the behavior of atomic nuclei in the degenerate electron gas: Do they move freely in it or do they oscillate near certain fixed points of equilibrium?

In the first case, we would deal with a mixture of the gases of electrons and nuclei. In the second case, we are faced with a crystal lattice which contains a free electron gas. In the latter case, the situation is similar to that in metals. The below-presented consideration will show that, at high densities, a version of the solid state is realized (see Ref. 264).

Let us consider the version with a crystal lattice and find the region of temperatures and densities, where this phase is stable. For the sake of simplicity, we consider identical nuclei.

We divide the medium into neutral spherical cells (Wigner–Seitz cells) with one nucleus and Z electrons in every cell. The cell radius is

$$R_c = \left(\frac{3}{4\pi n_A} \right)^{1/3} = \left(\frac{3\alpha Z m_p}{4\pi \rho} \right)^{1/3}, \quad (12.53)$$

where n_A is the density of nuclei, A is the mass number, Z is the nucleus charge, $\alpha = Z/A$, $\rho \approx n_A m_A \approx n_A A m_p$ is the mass density, m_A and m_p are, respectively, masses of a nucleus and a proton.

Assuming that the charge of the electron cloud in a cell is distributed uniformly, we find the potential energy of the nucleus as

$$U(r) = -\frac{3Z^2 e^2}{2R_c} + \frac{Z^2 e^2}{2R_c^3} r^2. \quad (12.54)$$

Here, r is a displacement of the nucleus from the cell center, $r < R_c$. The nucleus is located, thus, in a potential well with depth $U_0 = 3Z^2 e^2 / 2R_c$. The potential $U(r)$ is of the oscillatory type. By comparing it with the corresponding expression for an isotropic harmonic oscillator, we get the

frequency of oscillations as

$$\omega = \left(\frac{Ze^2}{\alpha m_p R_c^3} \right)^{1/2}. \quad (12.55)$$

The oscillator energy is $\varepsilon_n = \hbar\omega(n + 3/2)$, where $n = 0, 1, 2, \dots$ is the quantum number. At the zero temperature, $n = 0$. Therefore, the vibrational energy of the nucleus is

$$\varepsilon_0 = \frac{3}{2}\hbar\omega = \frac{3eh}{2\alpha m_p} \left(\frac{4\pi}{3}\rho \right)^{1/2}. \quad (12.56)$$

This value is significantly less than the depth of the potential well:

$$\frac{\varepsilon_0}{U_0} = \frac{2h}{e\sqrt{\alpha m_p R_A}} \frac{1}{Z} \left(\frac{R_A}{Z R_c} \right)^{1/2} \approx \frac{1}{Z} \left(\frac{R_A}{Z R_c} \right)^{1/2} \ll 1. \quad (12.57)$$

Here, R_A is the nucleus radius and, obviously, is less than R_c . This testifies to the strong coupling of a nucleus with the own cell and to the stability of the appropriate state.

Let us estimate the influence of temperature on a state of the system. The system of oscillators obeys the Bose–Einstein statistics (it is a system of bosons), and the properties of the system of oscillators in an equilibrium state are defined by the Planck distribution function

$$f_0(\varepsilon) = \frac{1}{\exp(\varepsilon/T) - 1}. \quad (12.58)$$

This yields the simple formula for the mean energy of oscillators, $\bar{\varepsilon} = \varepsilon_0 \text{cth}(\varepsilon_0/T)$. Then we can find the ratio of the mean energy and the depth of the potential well in a nucleus as

$$\begin{aligned} \frac{\bar{\varepsilon}}{U_0} &= \frac{\varepsilon_0}{U_0} \text{cth} \left(\frac{\varepsilon_0}{T} \right) \\ &= \frac{(4\pi/3)^{1/6} h}{e(\eta m_p)^{2/3}} Z^{-5/3} \rho^{1/6} \text{cth} \left[\frac{(3\pi)^{1/2} eh \rho^{1/2}}{T \eta m_p} \right]. \end{aligned} \quad (12.59)$$

Thus, the plasma phase in the nucleus arises at $\bar{\varepsilon} \gg U_0$. If $\bar{\varepsilon} \ll U_0$, the solid-state phase is realized.

Processes of Neutronization of a Substance. With increase in the substance density, numerous phase transitions occur. Consider the changes happening, first of all, in a substance completely composed by hydrogen,

because it is the most abundant substance in the universe, and the estimates for it are the simplest ones.

Consider a proton-electron plasma. Up to the values of density which are defined by the relation $\varepsilon_F \leq (m_n - m_p) c^2$, no phase transitions are observed. However, at $\varepsilon_F > (m_n - m_p) c^2$, the proton-electron gas becomes unstable relative to the creation of neutrons.

The neutronization of hydrogen and some astrophysical applications of this process are studied for a long time, beginning from work published in the 1930s (see Refs. 265–267).

This process becomes obvious if we consider two alternative states of the substance: the proton-electron “p + e” and neutron “n” ones.

Under ordinary conditions, a free neutron is unstable and decays into a proton, electron, and antineutrino. However, in the presence of a degenerate electron gas, a neutron can become a stable particle. Moreover, the process can run in the inverse direction. Due to the two-side processes of β -decay

$$n \rightarrow p + e + \bar{\nu}_e, \quad p + e \rightarrow n + \nu_e, \quad (12.60)$$

the thermodynamic equilibrium between the neutron and proton-electron phases is established in the substance. The condition of equilibrium has form of the relation between chemical potentials

$$\mu_p + \mu_e = \mu_n, \quad (12.61)$$

which is ordinary in thermodynamics. Here, μ_p, μ_e, μ_n are the chemical potentials proton, electron, and neutron, respectively. In the equality, we omit the chemical potential of a neutrino, because it equals zero due to its weak interaction with the medium (the interaction cross-section of a neutrino with electrons is of the order of 10^{-44} cm^2) and the zero rest mass.

Equation 12.61, with regard to the dependence of the chemical potentials on the density and temperature, yields the first equation for the determination of states of the substance. The second equation for the determination of thermodynamical quantities follows from the condition of electroneutrality

$$n_e = n_p. \quad (12.62)$$

Let us analyze the simplest case of a degenerate system. That is, we assume that all the components of the system have temperature below their degeneration temperature (less than the appropriate Fermi energies). Then we get the following formulas for chemical potentials:

$$\mu_e = \varepsilon_e = c \left(m_e^2 c^2 + \left(3\pi^2 \right)^{2/3} h^2 n_e^{2/3} \right)^{1/2}, \quad (12.63)$$

$$\mu_p = \varepsilon_p = m_p c^2 + (1/2 m_p) \left(3\pi^2 \right)^{2/3} h^2 n_p^{2/3}; \quad (12.64)$$

$$\mu_n = \varepsilon_n = m_n c^2 + (1/2 m_n) \left(3\pi^2 \right)^{2/3} h^2 n_n^{2/3}. \quad (12.65)$$

In view of these relations, Eq. 12.61 yields

$$n_n = \left[\frac{m_n}{m_p} n_p^{2/3} + \frac{2m_n}{(3\pi^2)^{2/3} h^2} \left(\varepsilon_e - (m_n - m_p) c^2 \right) \right]^{3/2}. \quad (12.66)$$

We introduce the notations $\alpha_* = \frac{m_n - m_p}{m_e} \approx 2.54$ and $n_* = \frac{m_e^3 c^3}{3\pi^2 h^3} = 5.87 \times 10^{29} \text{ cm}^{-3}$.

Then we can write

$$n_n = n_p \left[1 + 2 \frac{m_p}{m_e} \left(\frac{n_*}{n_p} \right)^{2/3} \left(\sqrt{1 + \left(\frac{n_p}{n_*} \right)^{2/3}} - \alpha_* \right) \right]^{3/2}. \quad (12.67)$$

By equating the expression in parentheses to zero, we obtain the estimate for the threshold of the creation of neutrons:

$$n_p^* = \frac{\alpha_*^2 - 1}{3\pi^2} \left(\frac{m_e c}{h} \right)^3 \approx 7.5 \times 10^{30} \text{ cm}^{-3}. \quad (12.68)$$

At densities $n_p > n_p^*$, neutrons become stable. These densities correspond to the limiting momenta of electrons $p_e^* \approx 2.2 m_e c$, and electrons are relativistic.

After overcoming the threshold, the density of neutrons grows very rapidly, which is seen in Fig. 12.7 presenting the dependence of the equilibrium ratio of the densities of neutrons and protons on the density of protons.

In the estimation, we use the nonrelativistic formulas for nucleons, which is true for densities less than the nuclear one. At densities significantly higher than ones, at which neutrons become stable, nucleons become relativistic, and the equilibrium densities of the components are equalized. In the ultrarelativistic region,

$$n_n = 8n_p. \quad (12.69)$$

The process of annihilation of electron-positron pairs and the inverse process (in the field of other particles) are possible. In this case, we deal with the energies of quanta in the interval

$$2 m_e c^2 \leq h\omega \leq h c n^{1/3}. \quad (12.70)$$

Pass now to the analysis of the state of equilibrium in the substance with nuclei with mass number A .

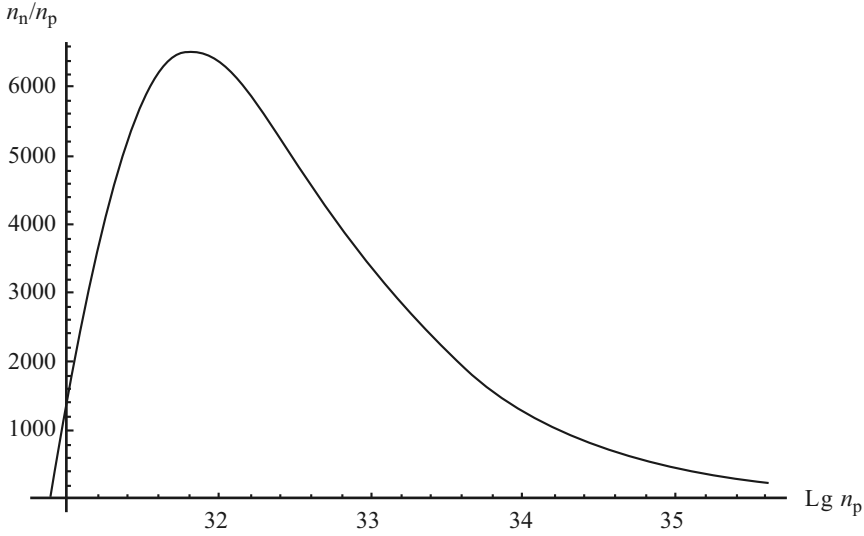


Fig. 12.7. Equilibrium ratio of the densities of neutrons and protons versus the density of protons.

As a basis model, we take the model of degenerate nucleon gas. According to this model, a nucleus is a potential well with width equal to its diameter and with depth of about 30 MeV, which is uniformly filled by neutrons and protons. The temperature of the nucleon gas can be taken, obviously, equal to zero, because the energy of excitation of nuclei is very large as compared to the thermal energy. Thus, nucleons in a nucleus form a completely degenerate Fermi-gas.

With regard to the potential well of a nucleus U_0 and the mean energy of the Coulomb interaction per proton, U_c , we have

$$\mu_p - U_0 + U_c + \mu_e = \mu_n - U_0. \quad (12.71)$$

The depths of the potential well for protons and neutrons is assumed to be the same, which is a consequence of the independence of nuclear forces on a charge state of a nucleon (a proton and a neutron are considered as different charge states of one particle named a nucleon). In a completely degenerate Fermi-gas, the chemical potential coincides with the limiting energy of particles (corresponding to the Fermi energy). Hence, the equilibrium between alternative states is defined by the relation

$$\varepsilon_p + U_c + \varepsilon_e = \varepsilon_n. \quad (12.72)$$

The densities of protons and neutrons in a nucleus are

$$n_p = \frac{3}{4\pi r_0^3} \frac{Z}{A}, \quad n_n = \frac{3}{4\pi r_0^3} (1 - Z/A), \quad (12.73)$$

where $r_0 = 1.3 \times 10^{-13}$ cm is a constant appearing in the formula $R_A = r_0 A^{1/3}$ for the nucleus radius. Here, it is assumed that the substance consists of one kind of nuclei, and, in addition, nucleons in a nucleus are distributed uniformly. For the limiting energies in the degenerate case, the following simple relations are valid:

$$\mu_e = \varepsilon_e = c \left[m_e^2 c^2 + (3\pi^2)^{2/3} h^2 n_e^{2/3} \right]^{1/2}, \quad (12.74)$$

$$\mu_p = \varepsilon_p = m_p c^2 + \left[(3\pi^2)^{2/3} h^2 / 2m_p \right] n_p^{2/3}, \quad (12.75)$$

$$\mu_n = \varepsilon_n = m_n c^2 + \frac{(3\pi^2)^{2/3} h^2}{2m_n} n_n^{2/3}. \quad (12.76)$$

Substituting the formulas for ε_p and ε_n in Eq. 12.72, we get

$$\begin{aligned} & (1 - Z/A)^{2/3} - (Z/A)^{2/3} \\ &= \frac{2m_p r_0^2}{h^2} \left(\frac{4}{9\pi} \right)^{2/3} \left(\varepsilon_e + U_c - \frac{m_n - m_p}{m_e} m_e c^2 \right). \end{aligned} \quad (12.77)$$

For medium nuclei, U_c is of the order of 1 MeV, and, at $n_e \leq 10^{31} \text{ cm}^{-3}$, the Fermi energy of electrons $\varepsilon_e \leq 1 \text{ MeV}$. Hence, at $n_e \leq 10^{31} \text{ cm}^{-3}$, the value of the right-hand side is close to zero. In this region of densities, the numbers of protons and neutrons are approximately equal with Z/A of the order of 0.5.

It is seen from Eq. 12.77 that the process of neutronization of protons in nuclei begins when the Fermi energy of electrons reaches the value

$$\varepsilon_e = \frac{m_n - m_p}{m_e} m_e c^2 - U_c. \quad (12.78)$$

Thus, beginning from densities $n_e = 10^{31} \text{ cm}^{-3}$, the increase in the density of electrons in nuclei is accompanied by the transformation of protons into neutrons. At the great values of ε_e , the ratio Z/A is a function of the density which decreases with increase in ε_e . In other words, at some value of ε_e , a nucleus becomes unstable relative to the reaction

$$(A, Z) + e \rightarrow (A, Z - 1) + \nu_e. \quad (12.79)$$

The limiting Fermi energy of electrons corresponding to this reaction is

$$\varepsilon_e = [M(A, Z - 1) - M(A, Z)] c^2. \quad (12.80)$$

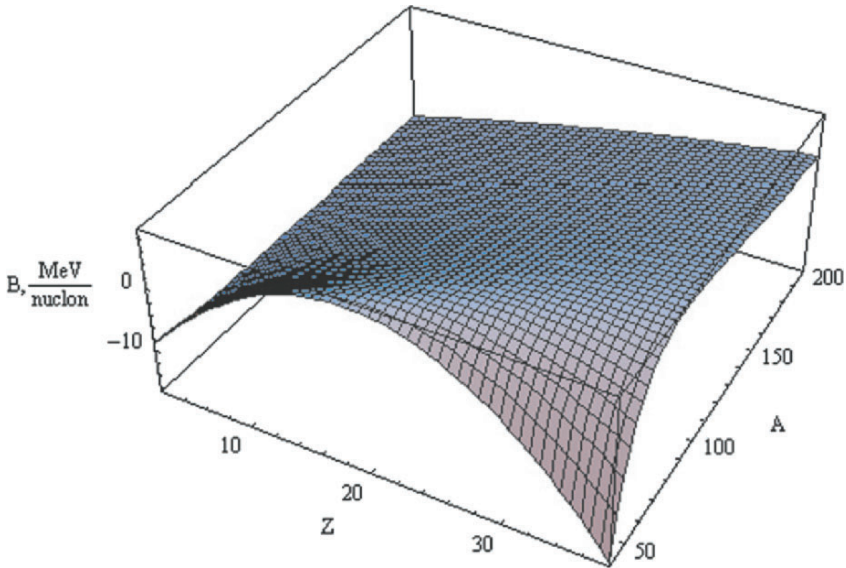


Fig. 12.8. Binding energy of nuclei versus their mass numbers and charges.

The nucleus mass can be written as

$$M(A, Z) = (A - Z)m_n + Zm_p + B(A, Z)/c^2, \quad (12.81)$$

where the binding energy as a function of the mass number and nucleus charge is given by the Weizsäcker formula (see Ref. 268):

$$B(A, Z) = c_0 A - c_1 A^{2/3} - c_2 \frac{Z^2}{A^{1/3}} - c_3 A (1 - 2Z/A)^2 + 34A^{-7/4} \delta, \quad (12.82)$$

where $c_0 = 15.7 \text{ MeV}$, $c_1 = 17.8 \text{ MeV}$, $c_2 = 0.71 \text{ MeV}$, $c_3 = 23.7 \text{ MeV}$, $\delta = 0, 1, -1$ for odd L , even L , and odd Z , respectively.

The Weizsäcker formula follows from the analysis of the Schrödinger many-particle equation in the Thomas–Fermi approximation. Numerical values of coefficients are derived to be sufficiently close to the empirical values.

In Fig. 12.8, we present the specific binding energy per nucleon versus the mass number A and nucleus charge Z .

Since the binding energy depends not only on the mass number, but also on the charge, Figs. 12.9 and 12.10 show, for the sake of clearness, the binding energy near the line of stability of nuclei $Z = f(A)$, whose form is discussed below. For the majority of nuclei at the middle of the Mendeleev periodic system of elements, the specific binding energy is of the order of 8 MeV/nucleon .

As seen from the figures, it is somewhat lesser for light nuclei due to a sharp increase in the Coulomb repulsion of protons at small distances and

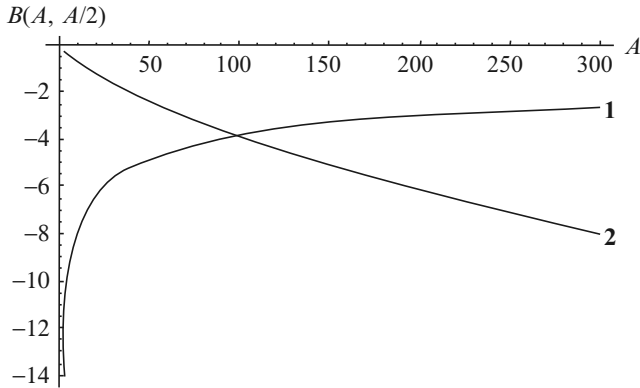


Fig. 12.9. Specific Coulomb (1) and surface Coulomb (2) binding energies versus the mass number near the line of stability of nuclei.

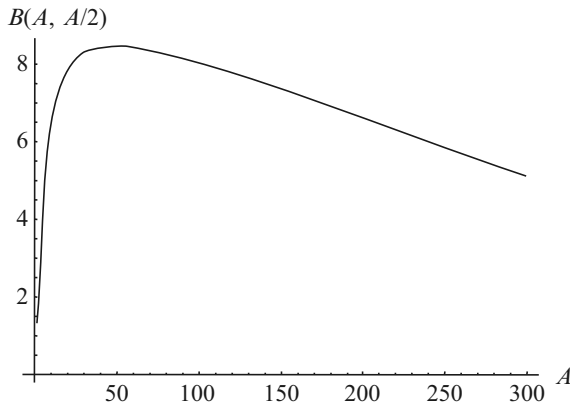


Fig. 12.10. Total binding energy versus the mass number near the line of stability of nuclei.

somewhat decreases for heavy nuclei due to a change in the relation between the surface and Coulomb energies.

Let us consider both cases of even and odd mass numbers.

- For odd A upon transition (Eq. 12.79), we get the stable nucleus $(A, Z - 1)$ up to a certain threshold value ε_e , at which the next transition to the nucleus $(A, Z - 2)$ occurs.
- For even A upon transition (Eq. 12.79) from even Z to odd $Z - 1$, we get an unstable nucleus. Therefore, the next transition to the stable nucleus $(A, Z - 2)$ occurs at once. This is conditioned by the last term in the Weizsäcker formula (Eq. 12.82).

With increase in the density, the number of protons in a nucleus decreases, and isotopes, being unstable under ordinary conditions, become

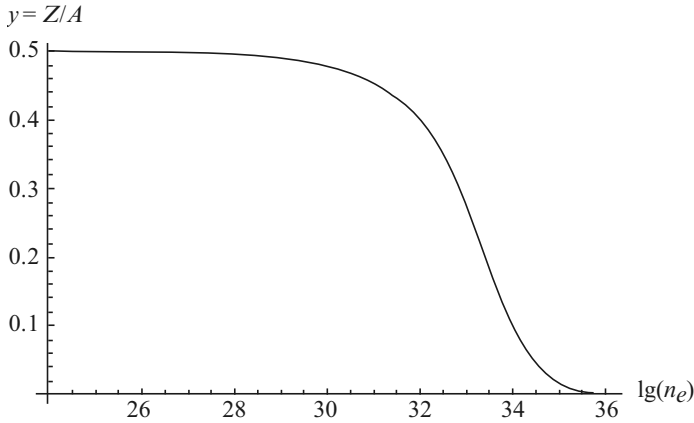


Fig. 12.11. Mean values of $y = Z/A$ versus the logarithm of the density of electrons.

stable at densities above a certain one. How long will the enrichment of nuclei by neutrons occur?

In the calculations of star's configurations, one uses the analytic formula derived from the estimate of the least value of Z/A at a given density of electrons n_e . This estimate can be easily derived by equating the binding energy to zero upon the fulfillment of conditions (Eq. 12.80) and the condition of neutrality $n_e = Zn_A$. As a result, we get the relation which can be approximated by the formula

$$\begin{aligned} A/Z = 2 + 0.01255 (p_F/m_e c) + 1.755 \cdot 10^{-5} (p_F/m_e c)^2 \\ + 1.376 \cdot 10^{-6} (p_F/m_e c)^3. \end{aligned} \quad (12.83)$$

This formula defines the dependence of the mean value of $y = Z/A$ on the density of electrons shown in Fig. 12.11.

The dependences of ratios Z/A for different nuclei (nuclei with different mass numbers A) on the density of electrons are taken into account in the calculations within star's models.

Eventually, at sufficiently high densities, the process of neutronization can lead to the formation of a nucleon gas, in which neutrons prevail. In addition, at sufficiently high densities of a substance, the reactions of synthesis can run even at zero temperature.

As a result of these reactions, the nuclear composition of the substance varies with increase in the density (from a substance with a single value of A , a mixture of nuclei with different A is derived).

Pycnonuclear Reactions and their Simplest Model. For the first time, the estimates of the probability of nuclear reactions running due to

an increase in the substance density were carried out in (see Refs. 269–271). These reactions caused by a great density are called pycnonuclear. The essence of these reactions is in that nuclei join at relatively small distances between them due to the tunnel effect.

To calculate the probability of this process, it is necessary to choose a model of potential barrier. We use the model of a Wigner–Seitz cell (Eq. 12.12). A nucleus is localized in a cell if the energy $(2\pi\hbar)^2/(m_A R^2)$ defined by the uncertainty relation for momenta is small as compared to the oscillator energy (Eq. 12.15). Using the ratio of these energies, we arrive at the inequality expressing the condition of localization of a nucleus:

$$\frac{6.6 \times 10^{-3}}{Z A^{2/3}} \rho^{1/6} \ll 1. \quad (12.84)$$

Consider a collision of two nuclei, being in adjacent cells with charges Z_1 and Z_2 ($Z_1 < Z_2$). The collision geometry is presented in Fig. 12.12.

The potential energy for the interaction of nuclei has the form

$$V(r) = \begin{cases} -\frac{3Z_1^2 e^2}{2R_1} + \frac{Z_1^2 e^2}{2R_1^3} r^2 + \frac{Z_1 Z_2 e^2}{s-r} - \frac{Z_1 Z_2 e^2}{|r-l|}, & r < \frac{R_1}{\cos \theta}, \\ -\frac{3Z_1 Z_2 e^2}{2R_2} + \frac{Z_1 Z_2 e^2}{2R_2^3} |r-l|^2 - \frac{Z_1^2 e^2}{r} + \frac{Z_1 Z_2 e^2}{s-r}, & r > \frac{R_1}{\cos \theta}. \end{cases} \quad (12.85)$$

The first two terms represent the interaction of the first nucleus with the electron cloud of the own cell, the third one corresponds to the interaction of two nuclei, and the last term corresponds to the interaction of

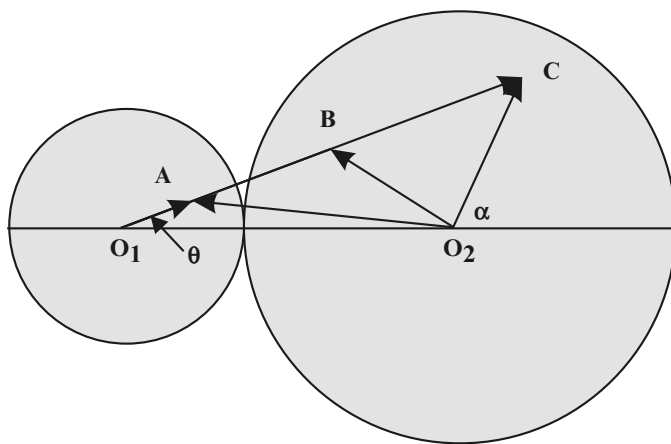


Fig. 12.12. Collision geometry of nuclei. $O_1O_2 = R_1 + R_2 = l$, $O_1C = S$, $O_2C = r'$, A , and B are the points on the way of the first nucleus to the point of the collision of two nuclei C .

the first nucleus with the electron cloud of the second cell. To estimate the probability, we simplify the potential.

We neglect the interaction of the nucleus with the second cell, while it is in the own cell (i.e., we neglect two terms in the upper row in Eq. 12.85). It is reasonable to replace the values of $|r - l| = O_2B$ and r in the lower row by their mean values, because, in the estimation of the probability, we should integrate over the second cell. We may take $|r - l| = 0.5R_2$ and $\bar{r} = R_1 + 0.5R_2$. Then the averaged potential takes the form

$$V(r) = \begin{cases} -\frac{3Z_1^2e^2}{2R_1} + \frac{Z_1^2e^2}{2R_1^3}r^2, & r < \frac{R_1}{\cos\theta}, \\ V_0 + \frac{Z_1Z_2e^2}{s-r}, & r > \frac{R_1}{\cos\theta}, \end{cases} \quad (12.86)$$

where $V_0 = -\frac{3Z_1Z_2e^2}{2R_2} + \frac{Z_1Z_2e^2}{2R_2^3}(0.5R_2)^2 - \frac{Z_1^2e^2}{\bar{r}}$.

The probability of the passage of the first nucleus from the own cell to the second one in the solid angle $\sin\theta d\theta$ is equal to $\frac{1}{2} \sin\theta d\theta \frac{\omega_1}{2\pi} \exp\left(-\frac{1}{\pi\hbar} \sqrt{2m_1} I(s, \theta)\right)$.

Here, m_1 and ω_1 are, respectively, the mass and frequency of oscillations of the first nucleus, and

$$I(s, \theta) = \int_{r_1}^{s-r_0} \sqrt{V(r) - E_1} dr, \quad (12.87)$$

where $E_1 = \frac{3}{2}\hbar\omega_1 - \frac{3}{2}\frac{Z_1e^2}{2R_1}$ is the energy of the ground state of the first oscillator, $r_1 = (3\hbar/Z_1e)^{1/2} (R_1^3/m_1)^{1/4}$ is the distance from the center of the first cell to the point, at which the energy level cross the barrier.

The probability of the presence of the second nucleus at a point of the volume element dV is $\psi_2^2(r')dV$, where $\psi_2(r') = a_2^{-3/2}\pi^{-3/4} \exp\left(-r'^2/2a_2^2\right)$ is the wave function of an oscillator in the ground state, $a_2 = \sqrt{3\hbar/m_2\omega_2}$ is the amplitude of oscillations of the second oscillator, m_2 and ω_2 are, respectively, the mass and frequency of these oscillations.

Using these relations, we get that this probability is equal to

$$a_2^{-3}\pi^{-3/2} \exp\left(-r'^2/a_2^2\right) r'^2 dr' \sin\alpha d\alpha. \quad (12.88)$$

The probability of the collision of two nuclei at any point C of the second cell and their fusion into one nucleus is

$$W = \frac{n\omega_1}{4\pi^{5/2}a_2^3} \int \exp\left[-\frac{r'^2}{a_2^2} - \frac{2\sqrt{2m_1}}{\hbar} I(s, \theta)\right] \times \delta\left(\theta - \arcsin \frac{r' \sin\alpha}{\sqrt{l^2 + r'^2 + 2lr' \cos\theta}}\right) r'^2 dr' d\Omega, \quad (12.89)$$

where $d\Omega = \sin\alpha \sin\theta d\alpha d\theta$, and $n \approx 27$ is the number of adjacent cells touching the first one.

As a result of the recent studies, the more exact and convenient relations for the probabilities of reactions were derived (see Ref. 272). The probability of the reactions (with tunneling through a Coulomb barrier) depends on the nucleus energy E in the following manner:

$$p(E) = \frac{\pi\sqrt{E_G/E}}{\text{Exp}\left(\pi\sqrt{E_G/E}\right) - 1}. \quad (12.90)$$

Here, E_G is the Gamow energy

$$E_G = 49.6 \cdot 10^{-3} (Z_1 Z_2)^2 (2\mu/m_p) \text{ MeV}. \quad (12.91)$$

An increase in the density of plasma leads to a modification of the probability and cross-section of the reaction. In this case, the principal point consists in that the probability of reactions with regard to the enhanced density

$$p(E) = \frac{\pi\sqrt{\frac{E_G}{E+E_s}}}{\text{Exp}\left(\pi\sqrt{\frac{E_G}{E+E_s}}\right) - 1}, \quad (12.92)$$

$$E_s = 0.144 \cdot 10^{-3} Z_1 Z_2 \left(\frac{10^{-9}}{D_s}\right) \text{ MeV}, \quad (12.93)$$

$$\frac{1}{D_s} = \frac{2}{\pi} \int_0^\infty dk \left(1 - \frac{1}{\varepsilon(0, k)}\right) \quad (12.94)$$

is not nullified, contrary to probability Eq. 12.90, at the zero and small energies of the nucleus, and the cross-section grows with energy:

$$\sigma(E) = \frac{S_0 (1 + a_1 E + a_2 E^2)}{\sqrt{E(E + E_s)} \left(\text{Exp}(\pi\sqrt{E_G/(E + E_s)}) - 1\right)}, \quad (12.95)$$

where $a_1 = 13.8$ and $a_2 = 1.246 \times 10^3$.

At high densities, the probabilities of reactions turn out sufficiently great, but the typical duration of reactions is of the order of tenths of millions of years for the characteristic parameters of massive stars.

We emphasize one more that the key role in our approach to the realization of nuclear processes is played by collective processes with regard to self-consistent electromagnetic fields arising in a target. In the following section, we will show that these fields modify the interaction potentials between nuclei and, as a result, increase the probability of the reaction.

In this case, of a special importance is the account of strong correlations, the disequilibrium in a nuclear system, and the appearance of the attraction between nuclei due to the polarization-related interaction in self-consistent fields.

Equilibrium in Neutron–Electron–Nucleus Plasma and the Phase Transition to a Nuclear Substance.

The reactions of synthesis run due to pycnonuclear reactions and lead to the appearance of nuclei with various values of A with increase in the density (see Ref. 273). Let us analyze the equilibrium composition of a nucleus, by using the Weizsäcker formula for the binding energy of nuclei at a fixed A :

$$\rho = \frac{n}{A} [(A - Z)m_n + Zm_p - B(A, Z)] + \frac{3}{4}c_e(Zn/A)^{1/3}, \quad (12.96)$$

where $c_e = (3\pi^2)^{1/3} 2\pi\hbar c = 6.1145 \times 10^{-11} \text{ MeV cm}$, n is the number of baryons in unit volume, $B(A, Z)$ is the binding energy of a nucleus. Here, the second term represents the energy of a degenerate relativistic electron gas.

We now find the energy density Eq. 12.96 at a fixed number of baryons.

The condition for this expression to be minimum is defined by the equations

$$\left(\frac{\partial \rho}{\partial Z}\right)_{n,A} = 0, \quad \left(\frac{\partial \rho}{\partial A}\right)_{n,Z} = 0. \quad (12.97)$$

We introduce a variable $y = Z/A$. Then Eq. 12.97 after the differentiation with respect to Z yields

$$\begin{aligned} \varepsilon_e &= a_e (ny)^{1/3}, \quad a_e = (3\pi^2)^{1/3} h, \\ a_e (ny)^{1/3} &= (m_n - m_p + 4c_3) - 2y(4c_3 + c_2 A^{2/3}). \end{aligned} \quad (12.98)$$

Equation 12.98 defines the dependence of $y = Z/A$ on the density n at a given A . The corresponding minimum energy density is

$$\begin{aligned} \rho(n, A) &= n \left[(m_n + c_3 - c_0) - \frac{1}{4}(m_n - m_p + 4c_3)y \right. \\ &\quad \left. + c_1 A^{-1/3} - \frac{1}{2}(4c_3 + c_2 A^{2/3})y^2 \right]. \end{aligned} \quad (12.99)$$

The second Eq. 12.97 yields

$$(m_n - m_p + 4c_3) - \frac{1}{3}c_1 A^{-1/3} - \frac{4}{3}c_2 A^{2/3}y^2 - 8c_3 y^2 - a_e n^{1/3}y^{4/3} = 0. \quad (12.100)$$

Jointly solving Eqs. 12.98 and 12.100, we get

$$y = \frac{Z}{A} = \left(\frac{c_1}{2c_2} \right)^{1/2} \frac{1}{\sqrt{A}} = \frac{3.54}{\sqrt{A}}. \quad (12.101)$$

In view of this solution and Eq. 12.98, we can derive the dependence of A and Z on the density of baryons in the most stable state of the degenerate electron-nucleus gas. It follows from the last result that, for every value of the density of baryons, there exists one stable nucleus with $A(n)$ and $Z(n)$. With decrease in the density, we get the limiting value $A = 56$. In this case, $Z/A \rightarrow 26/56$.

For the Fermi energy of electrons satisfying the inequality

$$Z\varepsilon_e \geq [Am_n - M(A, Z)]c^2, \quad (12.102)$$

a degenerate electron-nucleus plasma should contain free neutrons. A thermodynamic equilibrium is attained with the help of reactions

$$(A, Z) + Ze \rightleftharpoons An + Z\nu_e. \quad (12.103)$$

We get the condition for the appearance of free neutrons in a electron-nucleus plasma starting from the minimum of internal energy. With regard to the existence of free neutrons, we write the energy density as

$$\begin{aligned} \rho = & \frac{n - n_n}{A} [(A - Z)m_n + Zm_p - B(A, Z)] + nm_n \\ & + \frac{3a_e^2}{10m_n} n_n^{5/3} + \frac{3}{4} a_e [(Z/A)(n - n_n)]^{4/3}. \end{aligned} \quad (12.104)$$

Here, n_n is the density of neutrons, the penultimate term is the energy density of a nonrelativistic neutron gas, and the last term is the energy density of a degenerate relativistic electron gas. In Eq. 12.104, we took into account the condition of neutrality of the plasma, $Z(n - n_n) = A \cdot n_e$.

The condition for the thermodynamic equilibrium is the condition for the minimum of energy:

$$\left(\frac{\partial \rho}{\partial n_n} \right)_{n, A, Z} = 0, \quad \left(\frac{\partial \rho}{\partial Z} \right)_{n, n_n, A} = 0. \quad (12.105)$$

Equation 12.105 allows us to derive the conditions of equilibrium in the form of a system of equations:

$$\left(\frac{2m_n}{a_e^2} \right)^{3/2} \left[(c_3 - c_0) + \frac{1}{2} c_1 A^{-1/3} - \frac{2c_1 c_3}{c_2} \frac{1}{A} \right]^{3/2} = n_n, \quad (12.106)$$

$$\frac{1}{a_e^3} \left(\frac{2c_2}{c_1} \right)^{1/2} \left[(m_n - m_p + 4c_3) A^{1/6} - \sqrt{2c_1 c_2} A^{1/3} - 8c_3 \sqrt{\frac{c_1}{2c_2}} A^{-1/3} \right]^3 + n_n = n. \quad (12.107)$$

These relations allow us to calculate A and n_n for a given baryonic density n . An increase in the density can induce the full disintegration of nuclei conditioned by the increase in the Fermi energy of electrons. This threshold can be estimated from the requirement that the binding energy of nucleons be zero:

$$B(A, Z)/A = -(c_3 - c_0) - c_1 A^{-1/3} + 4c_3 y - (4c_3 + c_2 A^{2/3}) y^2. \quad (12.108)$$

This equation yields

$$y = \frac{2c_3 - \sqrt{4c_3^2 - (4c_3 + c_2 A^{2/3})(c_3 - c_0 + c_1 A^{-1/3})}}{4c_3 + c_2 A^{2/3}}. \quad (12.109)$$

The plot of this function is shown in Fig. 12.13. For the most probable (stable) nucleus, $y = \sqrt{c_1/2c_2} A^{-1/2}$. The plot of this function is also given in Fig. 12.13. The crossing of these plots defines the mass number of the last stable nucleus. As seen, prior to the very disintegration of the substance and the formation of a nuclear substance, the last stable nucleus has the mass number $A \approx 700$. Substituting this value in system Eqs. 12.106–12.107, we get that $n \approx 1.25 \times 10^{37} \text{ cm}^{-3}$ and $n_n \approx 7 \times 10^{36} \text{ cm}^{-3}$ on the threshold of the formation of the nuclear substance.

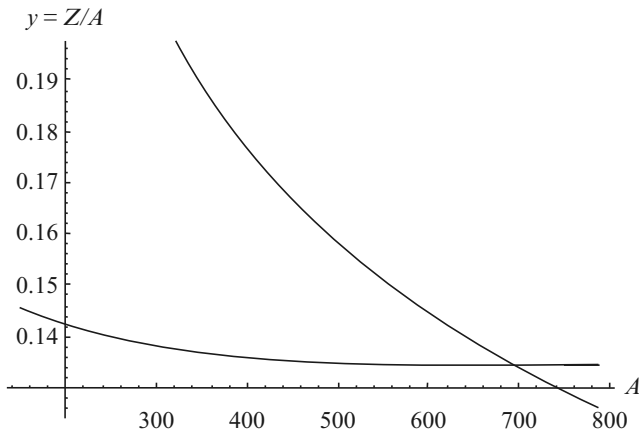


Fig. 12.13. Value $y = Z/A$ versus the mass number.

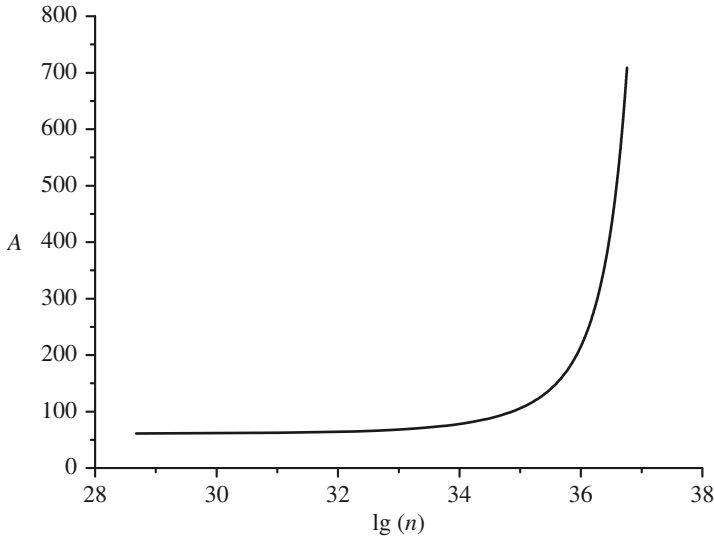


Fig. 12.14. Mass number A of the most stable nucleus versus the density of baryons (on a logarithmic scale).

Figure 12.14 demonstrates the dependence of the mass number of the most stable nucleus on the density of baryons up to the stability boundary on a logarithmic scale.

Thus, we have considered the properties of a neutral plasma consisting of free electrons and “bare” atomic nuclei at temperatures below the degeneration temperature of the electron gas. This nucleus-electron phase of the substance is formed at densities $\rho \geq 10\alpha^{-3}Z^4$, where Z is the ordinal number of a nucleus, and α is a number of the order of 1. The degeneration temperature of electrons in this plasma is sufficiently high. A substance in such a state is present in white dwarfs and the shells of baryonic stars (neutron-, quark-, and hyper-stars).

At $\varepsilon_e \geq 1 \text{ MeV}$ ($\rho \geq 10^7 \text{ g} \cdot \text{cm}^{-3}$), the neutronization of a substance occurs: due to the reactions of inverse β -decay, $(A, Z) + e \rightarrow (A, Z - 1) + \nu_e$, or positron-involved decays, $(A, Z) + e \rightarrow (A, Z - 1) + e^+ + \nu_e$, protons in nuclei are transformed in neutrons.

Both Z and the mass number A depend on ρ . The dependence of A on ρ is conditioned by the fact that light nuclei fuse in heavy ones due to the quantum-mechanical tunnel effect and disappear for this reason at high densities when the mean distances between particles become sufficiently small.

In this case, the parameters A and Z of atomic nuclei are functions of the limiting energy of electrons (or of the total energy density ρ).

When the limiting energy of electrons attains the value $Z\varepsilon_e = (Am_n - M)c^2$, where M is the mass of a nucleus with parameters A and Z , free neutrons appear in plasma. The nucleus–electron–neutron phase, which contains free neutrons along with atomic nuclei and electrons, is formed. At the threshold of the appearance of free neutrons,

$$\varepsilon_e \approx 23, \quad A/Z \approx 2.8, \quad \rho \approx 2.7 \cdot 10^{11} \text{ g} \cdot \text{cm}^{-3}. \quad (12.110)$$

Here, a free proton is unstable. Upon the further increase in the density, the ratio Z/A continues to decrease, and A grows. This phase is completed at $n \approx 10^{37} \text{ cm}^{-3}$ and $n_n \approx 7 \cdot 10^{36} \text{ cm}^{-3}$ (n_n is the density of free neutrons). Then the nuclear substance appears: namely, the electron-nucleon phase of the substance.

The mass number of the most stable nucleus, for which the internal energy is the least (the binding energy of nuclei is the greatest), depends on the density of baryons n in the nucleus-electron and nucleus-electron-neutron phases. With increase in n , it varies from $A = 56$ to $A \approx 122$ at the threshold of formation of the nucleus-electron-neutron phase and up to $A \approx 700$ prior to the full disintegration of atomic nuclei and the subsequent formation of the nuclear substance (see Eqs. 12.106–12.107).

Upon the study of states of the substance in the region of high density, the approximation of ideal gas is not correct. At densities of the order of the nuclear one, the important role is played by the nuclear forces of attraction. Their account leads, for example, to that Σ^- hyperons appear in the medium not at the density $\rho \approx 6.1 \times 10^{14} \text{ g cm}^{-3}$ (as follows from the approximation of ideal plasma), but at that by 3 times lesser, i.e., already at the density existing in the ordinary nuclear substance.

Besides electrons, leptons in a baryonic plasma are represented only by negative muons which appear at the nuclear density directly before the appearance of hyperons.

At sufficiently high densities when the medium contains almost all the hyperons and many resonances, only negative pions from the group of bosons can acquire stability.

The transition to the phase of a substance containing negative pions is realized when the limiting energy of electrons reaches the value $\varepsilon_e = m_\pi c^2$. In this phase, negative pions are condensed on the lowest energy level, i.e., all they are in the rest state. Just higher than the threshold of this phase, the medium is liberally filled by mesons, whose concentration becomes very soon comparable with the concentration of baryons. With the appearance of π^- -mesons, the densities of leptons are frozen. This happens for the reason identical to that for a system to have a π -meson in the rest state rather than an electron with the energy $\varepsilon_e > m_\pi c^2$, the former situation being energy-gainful.

When the mean distance between particles becomes of the order of $l \leq 2 \times 10^{-14}$ cm (it corresponds to the density $\rho > 6.1 \times 10^{40} \text{ g} \cdot \text{cm}^{-3}$), powerful attractive forces arise between baryons. The idea of a gas becomes completely wrong, and a baryonic fluid appears.

Equations of State of the Nuclear Substance in the Electron-Nucleus and Neutron-Electron-Nucleus Phases. Below, we write the equation of state for a substance in the electron-nucleus phase. In this case, the energy density is defined by the energy density of electrons and the density of nuclei n_k with charge Z_k and mass number A_k :

$$\rho = \left[m_n c^2 + \left(\frac{B(A_k, Z_k)}{A_k} \right) \right] \sum_k A_k n_k + \rho_e, \quad (12.111)$$

$$\sum_k A_k n_k = n \sum_k \frac{A_k}{Z_k} Z_k \frac{n_k}{n} = n \frac{\bar{A}}{Z} \sum_k Z_k \frac{A}{Z} = \sum_k \frac{A_k}{Z_k} \frac{n_k}{n}. \quad (12.112)$$

Since nuclei are carrying on zero oscillations around fixed points at sufficiently low temperatures, they do not contribute to the pressure. The last is defined by the pressure of degenerate gas electrons, $P \approx P_e$. In this case, the pressure can be described by Eq. 12.50. Consider two limiting cases.

- In the nonrelativistic case ($x_F \ll 1$), Eq. 12.50 yields $P \simeq \rho^{5/3}$.
- In the relativistic case ($x_F \gg 1$), we get $\rho \approx x_F^3$, $P \approx \rho^{4/3}$.

The equation of state is valid up to the value $x_F \approx 46$ corresponding to the energy density $\rho \approx 2.4 \times 10^{32} \text{ erg/cm}^3$. As was shown above, high energy densities are characterized by the appearance of the neutron component, for which the partial energy density and pressure [$x_n = p_n/m_n c = (3\pi^2)^{1/3} \hbar n_n^{1/3}/m_n c$ is the dimensionless limiting momentum for a system of neutrons] are

$$\rho_n = 4K_n \left[x_n (1 + 2x_n^2) (1 + x_n^2)^{1/2} - \ln \left(x_n + \sqrt{1 + x_n^2} \right) \right], \quad (12.113)$$

$$P_n = \frac{1}{3} K_n \left[x_n (2x_n^2 - 3) (1 + x_n^2)^{1/2} + 3 \ln \left(x_n + \sqrt{1 + x_n^2} \right) \right]. \quad (12.114)$$

The total energy density and pressure in this phase of the substance have the form

$$\rho = \rho_n + m_n c^2 A n_e / Z + \rho_e, \quad P = P_n + P_e \quad (12.115)$$

and are mainly defined by neutrons.

At present, the equation of state for the nuclear substance is intensively studied (see, e.g., Refs. 263, 274–282). In a wide range of pressures

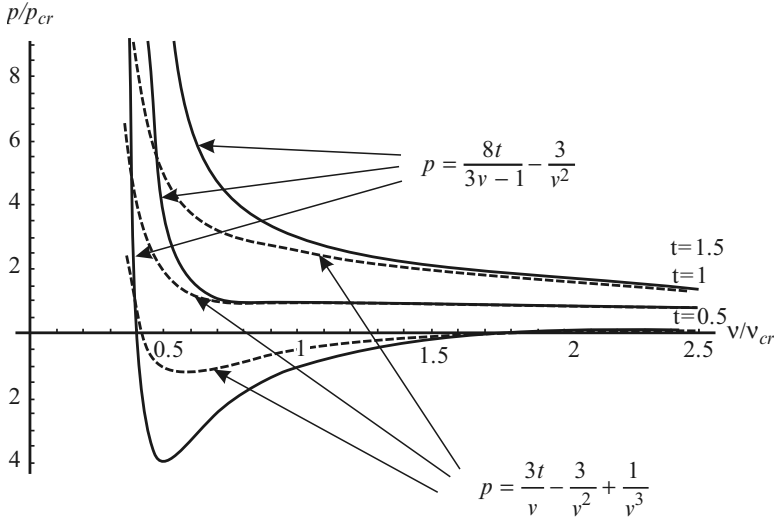


Fig. 12.15. Equations of state of an ordinary fluid and the nuclear substance in relative coordinates.

and temperatures, the nuclear substance is similar to the ordinary substance obeying the van der Waals equation of state. The equation of state looks most simply in the reduced coordinates $t = T/T_{cr}$, $p = P/P_{cr}$, and $v = V/V_{cr}$ ($T_{cr} = 17.3 \text{ MeV}$, $n_{cr} = 0.3 n_0$), where the variables are referred to those of the critical state. At the critical point, the difference between the liquid and gas phases disappears.

The van der Waals reduced equation of state takes the form

$$p = \frac{8t}{3v-1} - \frac{3}{v^2}. \quad (12.116)$$

For the nuclear substance, an analogous equation reads

$$p = \frac{3t}{v} - \frac{3}{v^2} + \frac{1}{v^3}. \quad (12.117)$$

In Fig. 12.15, we present the isotherms for the equations of state of a fluid and the nuclear substance in the reduced coordinates.

As seen, they differ each from other slightly, though they are referred to substances with different elementary components connected by forces of different nature at temperatures and pressures different in their scales. The reason for the similarity is related to that the forces have short-range character despite certain differences and are repulsive forces at very small distances and attractive forces at large distances.

At low temperatures, there exists a stable state of the baryonic substance, which corresponds to the substance of atomic nuclei in the liquid

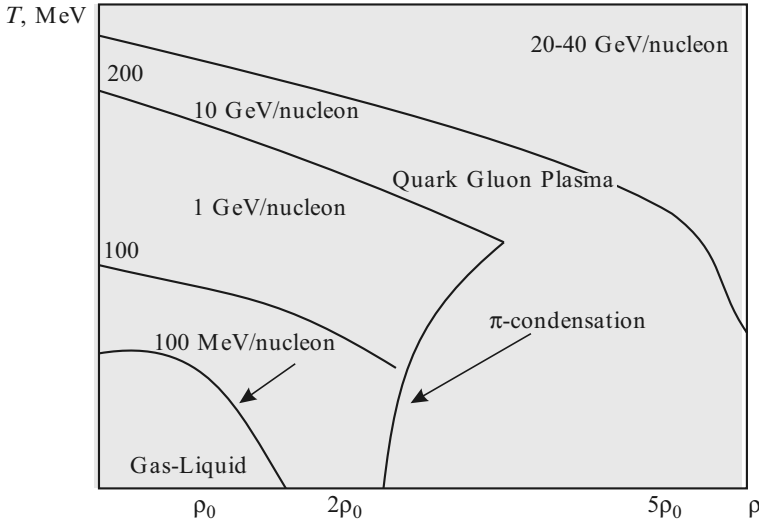


Fig. 12.16. Diagram of states of the nuclear substance.

state. When the temperature exceeds the critical value, the nuclear substance transits in the gas state, the nucleon gas (see Fig. 12.16).

At higher temperatures, the internal structure of nucleons is revealed. For example, isobars arise (in particular, hyperons which are composed from u , d , and heavier s quarks). Upon the further increase in the pressure, the appearance of quarks of the next generations (c , b , t) is possible until the phase transition of the baryonic substance in the quark-gluon phase (see Ref. 263).

At relatively low temperatures and with increase in the baryonic substance density, there occur the transition from the gas state in liquid one, then the transition to the phase with a pion condensate (see Ref. 283), and, finally, the transition in the “cold” quark-gluon phase.

Thus, we may expect phase transitions and the coexistence of the gas and liquid phases in the nuclear substance. In the nature, the possibility of a two-phase equilibrium in the nuclear substance can be realized at the bursts of supernovas and neutron stars. The experimental study of the baryonic substance is possible in a sufficiently wide range of densities and temperatures upon the creation of shock waves in the nuclear substance and in the presence of the processes of fragmentation of nuclei.

A two-phase system in a nucleus (the gas-fluid one) can arise in the evolutionary dynamical processes both from the side of high and low densities. In the first case, small bubbles of gas can arise in nuclear fluid. But if the substance evolves from the side of low densities, nucleons in the gas phase can form drops of the nuclear fluid (see Refs. 280–282, 284, 285).

From the surface of a dense nuclear system, there occurs the evaporation of nuclei with various atomic numbers. Such a fragmentation of big nuclei happens so that the whole spectrum of mass numbers is observed. The distributions of fragments over mass numbers were derived experimentally in the collisions of heavy nuclei (see, e.g., the experiments with colliding U^{238} nuclei [see Ref. 285]) and theoretically (see Refs. 280–286).

Experiments on the collision of heavy ions demonstrate the power distributions of fragments over mass numbers of the form $\approx A^{-2.65}$. Such spectra over mass numbers can be also explained within statistical models (see Refs. 287–288) with regard to the nonextensiveness (see Ref. 289) and within thermodynamic models (see Ref. 280) based on the theory of phase transitions (see Ref. 284). The comparison of the results derived in experiment and those following from the thermodynamic arguments leads to an interesting conclusion (see Ref. 280) about the formation of a fractal nuclear cluster with fractal dimensionality of its surface $D_f = 1.8$ inside a big nuclear system.

The fractal structure of the surface of a great nuclear system leads, naturally, to an increase in the surface energy $E_s(A, D_f)$ and a change of its dependence on the mass number. The latter should be taken into account in the formula for binding energy (Eq. 12.82) instead of the traditional term proportional to $E_s = -c_1 A^{2/3}$. Below, we present the estimates for the surface energy. For a dense cluster, the surface area $S = 4S_0$, where S_0 is the surface area where a cluster is positioned. In the general case (see Ref. 290), we get

$$S/S_0 = 4N^k, \quad (12.118)$$

where N is the number of particles in the cluster, and $0 < k < 1/3$. The surface area can be given directly in terms of the number of particles as

$$S = 4N^\gamma, \quad (12.119)$$

where $2/3 \leq \gamma \leq 1$. The quantities k and γ are connected with the fractal dimensionality of the cluster:

$$k = \gamma - 2/D_f. \quad (12.120)$$

Upon the formation of a cluster by means of the sequential addition of separate particles to it, the fractal dimensionality of a cluster can be determined from the condition of minimum of the free energy (see Ref. 291) and has the form

$$D_f = \frac{4D_\omega + d(2D_\omega - 4) + 5d^2}{5D_\omega - 4 + 5d}, \quad (12.121)$$

where d is the topological dimensionality of the space and D_ω is the fractal dimensionality of the trajectories of particles (in the case of the Brownian motion, $D_\omega = 2$). Using the above-presented relations, we get

$$E_s(A, D_f) = -c_1 A^{2/3} (A/A_0)^{\gamma-2/3}, \quad (12.122)$$

where A_0 is the mass number of elements participating in the formation of a cluster.

Bound structures analogous to atomic nuclei on the base of nucleons can arise also on the base of any baryons. In 1952, the first quasinuclear systems including Λ -hyperons besides nucleons were registered. These multibaryonic systems are called hypernuclei. At present, several tenths of different hypernuclei including Λ and Σ -hyperons are known. The lifetime of Λ -hyperons is of the order of 10^{-10} s, therefore the registration of such nuclei meets certain difficulties. These nuclei are mainly registered in nuclear photoemulsions as tracks created by nuclei up to their decay. There exist nuclei having no analogs among the ordinary nuclei such as ${}^4\text{H}_\Lambda$, ${}^5\text{He}_\Lambda$, and ${}^8\text{Be}_\Lambda$.

According to the modern ideas and experimental data, a stable nuclear substance possesses a certain nucleonic composition and the density n_0 . However, it is of interest to answer the question about the existence of atomic nuclei with other values of the density and binding energy.

An attempt to answer this question was made in the 1960s by Migdal (see Refs. 283, 292–293). He developed the theory of electron and pion condensation in superdense nuclei at the expense of the creation of electrons and pions from vacuum. In theory, the pion condensate can appear at the nuclear density $2 \div 3n_0$. Since this question is very significant for the understanding of the processes induced by REBs in the experiments performed at the Electrodynamics Laboratory “Proton-21”, we will discuss the main conclusions made by Migdal.

12.1.4. Electron and Pion Condensations in Nuclei: Anomalous Nuclei and Other Exotic Nuclear States

In physical systems with electric potential exceeding the rest energy of an electron–positron pair, vacuum is polarized. In sufficiently strong electric fields, the system accumulates electrons created from vacuum. In this case, the electrons screen the system charge and decrease the Coulomb energy of the system admitting no existence of systems with high Coulomb energy.

The state of electrons, like that of any other fermions (particles with half-integer spin), is described by the Dirac relativistic equation (see Ref. 294) which was deduced for the quantum-mechanical description of fermions with regard to relativistic properties. Like to the derivation of the

Schrödinger equation, Dirac started from the formula for the Hamilton relativistic function depending on the particle momentum:

$$H = \sqrt{p^2 c^2 + m^2 c^4}. \quad (12.123)$$

We recall that the relativistically covariant forms of a momentum p^μ and a coordinate x_μ reads

$$p^\mu = (p^0, p^1, p^2, p^3) = (E/c, p_x, p_y, p_z), \quad (12.124)$$

$$p_\mu = (p^0, -\mathbf{p}), \quad (12.125)$$

$$p_\mu p^\mu = E^2/c^2 - \mathbf{p} \cdot \mathbf{p} = m^2 c^2, \quad (12.126)$$

$$x_\mu = (ct, -\mathbf{x}), \quad (12.127)$$

$$x^\mu = (ct, \mathbf{x}), \quad (12.128)$$

$$ds^2 = g_{\mu\nu} dx^\mu dx^\nu = dx^\mu dx_\mu. \quad (12.129)$$

In this case, the transition to operators is realized by the relations

$$\nabla^\mu = \frac{\partial}{\partial x_\mu}, \quad (12.130)$$

$$H \rightarrow i\hbar \frac{\partial}{\partial t}, \quad (12.131)$$

$$p^\mu \rightarrow i\hbar \frac{\partial}{\partial x_\mu} \equiv i\hbar \nabla^\mu, \quad (12.132)$$

$$p_\mu \rightarrow -i\hbar \frac{\partial}{\partial x^\mu}. \quad (12.133)$$

For a four-component wave function $\psi = \begin{pmatrix} \psi_1 \\ \psi_2 \\ \psi_3 \\ \psi_4 \end{pmatrix}$, the relativistic equation for an electron reads as follows:

$$\left[i\hbar \left(\gamma^\mu \frac{\partial}{\partial x_\mu} \right) - mc \right] \psi = 0. \quad (12.134)$$

Here, $\gamma^0 = \beta$, $\gamma^i = \beta \alpha_i$, $\gamma^0 = \begin{pmatrix} I & 0 \\ 0 & -I \end{pmatrix}$, $\gamma^i = \begin{pmatrix} 0 & \sigma_i \\ -\sigma_i & 0 \end{pmatrix}$,

$$\beta = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \quad \sigma_1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \sigma_2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix},$$

$$\sigma_3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \quad \gamma^\mu \frac{\partial}{\partial x_\mu} = \frac{\gamma^0}{c} \frac{\partial}{\partial t} + \gamma \nabla.$$

For an electron in an electromagnetic field, on expressing $\mathbf{E} = -\frac{1}{c}\frac{\partial \mathbf{A}}{\partial t} - \nabla\varphi$ with the help of potentials $A^\mu = (\varphi, \mathbf{A})$, the Dirac equation is reduced to

$$\left(\widehat{p} - e\widehat{A}/c - mc\right)\psi = 0, \quad (12.135)$$

where $\widehat{A} = \gamma^\mu A_\mu = g_{\mu\nu}\gamma^\mu A^\nu = \gamma^0 A^0 - \gamma\mathbf{A}$ and $\widehat{p} = \gamma^\mu \partial/\partial x^\mu$.

In the analysis of the Dirac equations, it is convenient to introduce the dimensionless variables $p_0 = p_F/m_e c$ and $\omega = \varepsilon/m_e c^2$. Then

$$p_0 = \sqrt{\omega^2 - 1}, \quad (12.136)$$

$$n_e \lambda_c^3 = 3\pi^2, \left(\omega^2 - 1\right)^{3/2}, \quad (12.137)$$

where $\lambda_c = h/m_e c$ is the Compton wavelength of an electron, $\lambda_c = 2.426 \cdot 10^{-10}$ cm. We will measure the coordinates in units of the Compton wavelength of electrons and time in units of $t^* = \hbar/m_e c^2 = 1.288 \times 10^{-21}$ s.

As usual, we consider firstly the motion of an electron in the centrally symmetric Coulomb field of a nucleus described by the vector potential

$$A_k = (-eZ/r, 0, 0, 0). \quad (12.138)$$

Upon the motion in such a field, the momentum and parity are conserved. In this case, a solution of the equation can be represented in the standard form (Ref. 293)

$$\psi = \begin{pmatrix} \phi \\ \chi \end{pmatrix} = \begin{pmatrix} f\Omega_{jlM} \\ (-1)^{\frac{1+l-l'}{2}} g\Omega_{j'l'M} \end{pmatrix}, \quad l = j \pm 1/2, \quad l' = 2j - l. \quad (12.139)$$

For radial functions $F(r) = rf_\omega(r)$, $G(r) = rg_\omega(r)$, the Dirac equation is reduced to the system of equations

$$\frac{dG}{dr} + \frac{\kappa}{r}G - (\omega + 1 - V)F = 0, \quad (12.140)$$

$$\frac{dF}{dr} - \frac{\kappa}{r}F + (\omega - 1 - V)G = 0, \quad (12.141)$$

where $\kappa = \mp(j + 1/2)$. The upper and lower signs are taken synchronously with those in the formula for l .

From this system, we can get one second-order equation for the function G , by excluding F :

$$G'' + \frac{V'}{1 + \omega - V} \left(G' + \frac{\kappa}{r}G \right) + \left[(\omega - V)^2 - 1 - \frac{\kappa(\kappa + 1)}{r^2} \right] G = 0. \quad (12.142)$$

After the substitution

$$G = \sqrt{1 + \omega - V} u(r), \quad (12.143)$$

the foregoing equation is reduced to the Schrödinger equation

$$\frac{d^2 u(r)}{dr^2} + k^2(r) u(r) = 0, \quad (12.144)$$

where $k^2 = 2(E - U(r))$, $E = \frac{\omega^2 - 1}{2}$, $U(r) = \omega V - \frac{1}{2}V^2 + \frac{\kappa(\kappa+1)}{2r^2} + U_s$.

The first terms represent the effective potential of the Klein–Gordon–Fock equation, and the term U_s is related to spin effects:

$$U_s = \frac{1}{4} \left[\frac{V''}{1 + \omega - V} + \frac{3}{2} \left(\frac{V'}{1 + \omega - V} \right)^2 - \frac{2\kappa V}{r(1 + \omega - V)} \right]. \quad (12.145)$$

The derived equation differs significantly from the Schrödinger equation in a potential field, because the effective potential depends on energy and external electromagnetic fields. But near the energies with $\omega = -1$, the dependence on energy is weak. Under the assumption of the smallness of external fields, we can use the approximation

$$U(r) = \frac{\xi}{r} - \frac{\xi^2 - \kappa^2 + \frac{1}{4}}{2r^2}. \quad (12.146)$$

Consider the quasiclassical solutions of the Dirac equation. For this purpose, we introduce the amplitudes and phases of wave functions as

$$G = a \exp(iS), \quad F = b \exp(iS). \quad (12.147)$$

Here, as in the well-known quasiclassical WKB approximation (Ref. 295), the amplitudes a and b depend slightly on coordinates as compared to the phase S .

The density of electrons is defined, as usual, by the sum of the moduli squared of the electron wave functions over all one-particle states:

$$n_e = \frac{1}{3\pi^2 \hbar^3} \left(\frac{\varepsilon_F^2}{c^2} - m_e^2 c^2 + \frac{e^2}{c^2} V^2 - 2 \frac{e\varepsilon_F}{c^2} V \right)^{3/2}. \quad (12.148)$$

With regard to the relation between the electric potential φ and the potential energy $-e\varphi = V$, we get the equation for the self-consistent potential in a nucleus

$$\Delta V = -4\pi e^2 [n_e(r) - n_p(r)], \quad (12.149)$$

where the distribution of protons in the nucleus $n_p(r)$ is assumed to be set. Using the density of electrons in terms of the potential energy with regard to the central symmetry, we obtain the equation for V :

$$\frac{1}{r^2} \frac{d}{dr} \left(r^2 \varepsilon \frac{dV}{dr} \right) = -4\pi e^2 \left[\frac{1}{3\pi^2} (\omega^2 - 1 + V^2 - 2\omega V)^{3/2} - n_p(r) \right]. \quad (12.150)$$

Now consider the main positions of the theory of a vacuum electron shell of nuclei with great charge. In the region of energies with $\omega \approx -1$, we can get an approximation of the Thomas–Fermi relativistic equation for the calculation of the potential and, hence, the electron density in a nucleus as

$$\Delta V = -4\pi e^2 \left[\frac{1}{3\pi^2} (V^2 + 2V)^{3/2} - n_p(r) \right]. \quad (12.151)$$

It follows from this formula that the density of electrons is different from zero only in a bounded region of the space, $r < r_a$, in which the condition $V(r) < -2$ is satisfied. The shell boundary is defined by the condition $V(r_a) = -r_a V'(r_a) = -2$.

The distribution of the electric field in a nucleus calculated according to Eq. 12.151 for various charges of nuclei is presented in Fig. 12.17.

For great charges, the field turns out to be concentrated near the nucleus boundary. For the external observer, the self-consistent charge of a nucleus is considerably decreased, and the contribution of the electromagnetic component to the nucleus energy is mainly given by the kinetic energy of the electron vacuum shell and the surface energy. Thereby, the Weizsäcker formula for nuclei with a great number of protons is essentially changed. It is

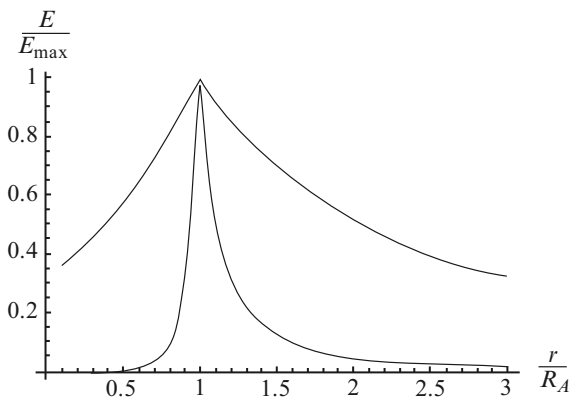


Fig. 12.17. Electric field in a nucleus versus the distance from the center (the field is referred to the field on the nucleus boundary, and the distance is referred to the nucleus radius). The lower and upper curves correspond, respectively, to $A = 1500$ and $A = 500$.

worth noting that the still greater changes in the properties of nuclei occur under the pion condensation.

The main sense of the idea consists in the following. With increase in the density of the nuclear substance ($\rho > \rho_c \geq \rho_0$), The loss of the baryonic subsystem in energy can be compensated by a gain at the expense of the appearance of a π -condensate. In this case, the system is compressed until it passes in a new equilibrium state with positive compressibility. These new states can correspond to anomalous nuclei, the thought about the possibility of their existence being advanced in (see Ref. 283).

The answer to the question whether metastable or even stable states at high densities $\rho > \rho_c$ exist is defined by the competition of two large numbers: the positive energy of the baryonic subsystem and the negative one of the π -condensate. Moreover, in the region of high densities, both values are not exactly known. As for the baryonic subsystem, there exist many equations of state which describe well the nuclear data (equilibrium density, binding energy, and symmetry-related energy), but reveal the different dependences on the density and different compressibilities. These models give different results for the energy of a baryonic system. The π -condensate energy can also be determined only with the help of model-dependent estimates. Bound anomalous states arise upon the use of the comparatively soft equations of state for the nucleonic subsystem in combination with the models of a sufficiently developed π -condensate.

It turns out that pion “drops” can arise in the nuclear substance due to the strong πN -interaction depending on the ratio N/Z . In the region of localization of pions of the same kind, it is energetically profitable to violate the isotopic invariance of the nuclear substance. As a result, the region of localization is enriched by nucleons of the same kind.

The lifetime of such drops can be large. Indeed, only neutrons are positioned near π^- mesons in the region of localization. Therefore, pions are absorbed only with the drop surface, and the lifetime of these objects is remarkably more than that of a nucleus. It follows from estimates that the number of pions in a drop $Q \approx 50$ for a drop of $a \approx 5$ fm in size, and the binding energy per pion is ~ 2 MeV. These estimates depend significantly, like in the case of anomalous nuclei, on the force of the πN -interaction and on the used equation of state.

The predictions become more optimistic if the extraneous hadrons are kaons. For example, already several K^+ -mesons is sufficient in order to create a K^+ -bubble (see Refs. 296–297), inside which nucleons are absent. The discovery of such objects would give an additional information about both the hadron-nucleus interaction and the equation of state of the nuclear substance at densities strongly different from the equilibrium one.

Let us analyze the possibility for stable superdense and superheavy nuclei to exist. It follows from the theory of pion condensation that there exist two possible regions of stability of anomalous nuclei: with superdense ($Z \approx N, Z \leq 10^2$) and superheavy ($Z \approx N, Z \geq 10^3$; in this case, the electric charge of baryons is almost completely screened by the π -condensate and electrons) nuclei (see Refs. 298–299). The stability region of superdense nuclei, like ordinary ones, is bounded by the possibility of their fission.

In superdense nuclei, like ordinary ones, a great role is played by exchange interactions. Therefore, $Z_h \approx N$, where Z_h is the total electric charge of baryons and π -condensate mesons. Too light superdense nuclei cannot exist due to a large contribution of the positive surface energy. In addition, no π -condensate appears in a low-nucleon system. For $A > 200, \dots, 400$, the Coulomb energy enters the play, and superdense nuclei terminate to be stable relative to the fission.

Nuclei become again stable relative to the fission at $Z \gg 1/e^3 \approx 1600$. Indeed, in the limiting case with $Z \gg 1/e^3$, the interior of a superheavy nucleus is an electroneutral plasma of baryons, π -condensate mesons, electrons, and negative muons. We emphasize that there is no limitation from above on A for such nuclei, and, in principle, nuclei-stars can exist (see Refs. 300–302).

Anomalous nuclei were experimentally searched for in natural specimens, among the products of the interaction of high-energy particles with a substance, in the products of fission, in collisions of heavy ions, etc. (see Refs. 303–305).

In nature, anomalous nuclei can be formed only under extreme conditions, for example, inside neutron stars upon their formation in the bursts of supernovas. In this case, the interiors of some neutron stars can be huge bunches of the anomalous substance. In natural specimens on the Earth, they can apparently be only in negligible amounts. Under laboratory conditions, superdense nuclei would be theoretically derived in collisions of high-energy heavy ions which, however, will meet the hampering obstacles such as the strong heating ($T \approx m_\pi$) and, possibly, the insufficiently high density [$\rho \sim (2 \dots 3)\rho_0$] of nuclear fireballs.

Besides anomalous nuclei, other exotic objects diverse in their nature can be observed. They are the long-lived localized states of hadrons (of pions and kaons) in the form of drops and bubbles in the nuclear substance (see Refs. 296–297), strange anomalous nuclei and stars (see Refs. 306–309), σ -nuclei (see Refs. 310–312), etc.

An increase in the substance density is associated in any case with increase in the electric field intensity, since the proton of a nucleus turn out to be concentrated in a small volume. At high values of the electric field,

the processes of creation of particles from vacuum can occur. In the case of systems with electric potential strongly exceeding the energy corresponding to the rest mass of electrons, the electrons arising from vacuum screen the system charge and decrease the Coulomb energy by enhancing the system stability.

The processes of creation of particles from vacuum involve not only electrons, but other particles as well. It is obvious that the most significant are the processes with the participation of the lightest (after electrons) particles. First of all, such particles are π -mesons. The importance of processes with their participation is related, in addition, with the fact that they as bosons can be accumulated in the ground state in the amount which is not limited by the quantum statistics.

The charge distribution for a vacuum electron shell of a nucleus with large charge is well described in the approximation based on the Thomas–Fermi relativistic equation used in the calculation of the distribution of the electric field potential in a nucleus. At $Ae^3 \ll 1$, electrons and negative muons play an insignificant role, but essential is the inhomogeneity of the pion charge distribution in a nucleus. At $Ae^3 > 1$, the screening by electrons becomes significant. With a further increase in A , the screening by negative muons should be taken into account as well. At the density $\rho^* \approx 10\rho_0$, a half of the positive charge is compensated by fermions (e^- , μ^-), and the other half is neutralized by π -mesons of the condensate. At $\rho \ll \rho^*$, a greater contribution to the screening inside a nucleus is introduced by π -mesons, and the field outside a nucleus is screened by electrons.

Thus, in the limiting case of a big superheavy nucleus ($Ae^3 \gg (\rho^*/\rho)^2$), it contains the electroneutral plasma including barionic quasiparticles, condensed pions, electrons, and negative muons in its interior. It is clear that the energy of such nuclear substance is less, in this case, than the energy of a pure neutron substance. Therefore, the account of both the accumulation of fermions and the nonuniformity of the pion charge distribution improves the conditions of stability of anomalous nuclei, the values of nuclear constants being taken the same as in the case without the consideration of these effects.

The Coulomb energy of superheavy nuclei is reduced to the surface energy (see Fig. 12.16). Therefore, the stability region for superheavy nuclei begins from $A \geq 10^3$, rather than from $A \geq 10^5$, as it would be without the account of the indicated effects. The observed charge of a superheavy nucleus turns out to be $\sim A^{1/3}$. These changes significantly modify the Weizsäcker formula for the binding energy.

In the next sections, we will show that the development of the ideas of Migdal on the basis of the conception of electrodynamic pycnonuclear

nucleosynthesis proposed in this work allows us to say about the formation of stable superheavy nuclei under the nonequilibrium conditions realized in the experiments performed at the Electrodynamics Laboratory “Proton-21”.

Reference 209 and Chapter 3 of this book contain the results of experiments which can be interpreted, according to the mentioned conception, as experimental observations of anomalous nuclei.

12.1.5. Nonequilibrium Thermodynamic Relations for Many-Particle Systems

The presence of energy sources in the system leads to the appearance of universal power asymptotics of the distribution function over energy for a wide class of interaction potentials between particles and forms of the collision integral.

The theoretical investigations of this phenomenon were first carried out in Ref. 313. The nonequilibrium distribution functions of electrons over energy in metals were measured, and it was shown that they have sections with power asymptotics $f(\varepsilon) \propto A\varepsilon^{s_\varepsilon}$ (see Ref. 314).

It was shown that deviations from the equilibrium are defined by the presence of flows in the phase space. The quasistationary nonequilibrium states of particles in such systems are similar to the Kolmogorov spectra of waves in the turbulent state (see Ref. 315). The power distributions of particles are derived as the stationary solutions of kinetic equations and nullify the collision integral.

Tsallis (see Ref. 289) generalized the Boltzmann–Gibbs traditional thermostatics to the states with the essential correlations and interaction between particles, which leads to the nonextensiveness of entropy and to quasipower distribution functions.

The formalism of Tsallis consists in the formal change of exponential and logarithmic functions in the relations of statistics and thermodynamics by their generalizations expressed through power functions:

$$\begin{aligned}\ln(x) &\rightarrow \ln_q(x) = \frac{x^{1-q} - 1}{1 - q}, \\ \exp(x) &\rightarrow \exp_q(x) = (1 + (1 - q)x)^{1/(1-q)},\end{aligned}\tag{12.152}$$

with some numerical parameter q . We note that, as q tends to 1, $\ln_q(x)$ and $\exp_q(x)$ pass in the ordinary logarithm and exponent. It is easy to show that the given definitions yield directly the relation

$$\ln_q(A + B) = \ln_q(A) + \ln_q(B) + (1 - q) \ln_q(A) \ln_q(B).\tag{12.153}$$

According to this relation, the new form of entropy (q -entropy) introduced in Ref. 289,

$$S_q = - \sum_i p_i^q \ln_q(p_i) = (1 - \sum_i (p_i^q)/(q-1)), \quad (12.154)$$

is not already an extensive function. If the whole system is divided into two independent subsystems A and B , then we get

$$S_q(A+B) = S_q(A) + S_q(B) + (1-q)S_q(A)S_q(B). \quad (12.155)$$

It follows from Eq. 12.155 that the parameter q is the measure of nonextensiveness of a system. As seen, the quantity q is limited by nothing and can take any real values. However, some limitations can appear in various specific problems. If $q > 1$, then the new entropy (q -entropy) takes the Boltzmann standard form.

In Eq. 12.154, p_i is the probability of the i -th state of the system and satisfies the normalization condition

$$\sum_i p_i = 1. \quad (12.156)$$

It is convenient to introduce the q -averaging of observables:

$$\langle O \rangle = \frac{1}{Z_q} \sum_i O_i p_i^q. \quad (12.157)$$

Here Z_q is the normalizing factor defined by the relation

$$Z_q = \sum_i p_i^q. \quad (12.158)$$

It is seen that the role of the probability distribution in the calculation of mean physical values is played by the q -distribution

$$\bar{p}_i = p_i^q / Z_q. \quad (12.159)$$

The equilibrium distribution p_i is usually determined from the entropy maximum condition upon the satisfaction of additional restricting conditions. Besides the condition of normalization for the distribution function, such conditions are the requirements for various physical quantities, for example for the energy E and charge Q of the system, to be fixed.

As known, the simplest mean to find the extremum of a function with certain limitations is the introduction of the appropriate Lagrange multipliers. In our case, this leads to the variational principle

$$\delta S + \alpha \sum_i \delta p_i - \beta_T \delta E + \gamma \delta Q = 0, \quad (12.160)$$

where $E = \sum_i E_i \bar{p}_i$, $Q = \sum_i Q_i \bar{p}_i$, and E_i and Q_i are, respectively, the energy and charge of the i -th state. The quantities α , γ , and β_T are the appropriate Lagrange multipliers. The multiplier β_T is connected, as usual, with temperature:

$$\beta_T = \frac{1}{T} = \left(\frac{\partial S}{\partial E} \right) \Big|_{Q,V}, \quad (12.161)$$

and the multiplier γ depends on the chemical potential μ :

$$\gamma = \frac{\mu}{T} = \left(\frac{\partial S}{\partial Q} \right) \Big|_{E,V}. \quad (12.162)$$

The variational equation Eq. 12.160 yields the power-like distribution

$$\bar{p}_i = \frac{1}{Z_q} \left[\exp_q (-\beta (E_i - \mu Q_i)) \right]^q, \quad (12.163)$$

where the statistical sum (the normalizing factor) is defined by the relation

$$Z_q(\beta, \mu, V) = \sum_i \left[\exp_q (-\beta (E_i - \mu Q_i)) \right]^q. \quad (12.164)$$

The quantity β is connected with temperature by the relation

$$T = \frac{\beta^{-1} + (q-1)(E - \mu Q)}{1 + (1-q)S}. \quad (12.165)$$

It was shown in (see Ref. 287) that the quantity

$$\bar{T} = \beta^{-1} + (q-1)(E - \mu Q) \quad (12.166)$$

is most adequate to the notion of physical temperature in nonextensive states. The above relations are especially important in the analysis of the process of evaporation of big nuclei. Below, we present the results of applications of the nonextensive statistics to the determination of the equation of state for nucleons in a nucleus (see Ref. 316).

The necessity to use the nonextensive statistics for the nuclear substance with extremely high densities is related to the enhancement of the role of the effects of memory, large-scale interaction, and non-Markovian processes in the course of the establishment of a stationary state.

In the estimates, we use the relativistic nonlinear model of a many-particle system consisting of interacting nucleons and mesons (see Ref. 317).

In Fig. 12.18, we present the results of calculations of the equation of state of the nuclear substance as a function of the parameter of nonextensiveness q .

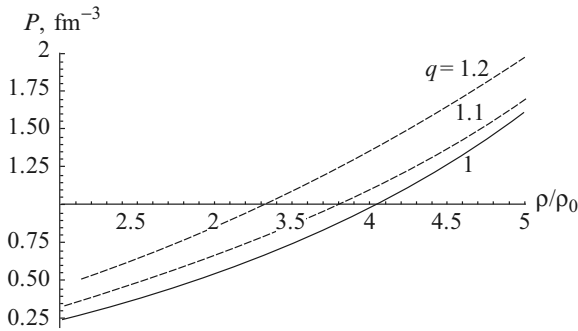


Fig. 12.18. Equation of state for hadrons with regard to the nonextensiveness of a system. The dependence of pressure on the nuclear density for various values of q .

Changes in the equation of state induce changes in the contribution of the strong interaction to the binding energy. This contribution turns out dependent on the parameter of nonextensiveness q which is defined by correlations present in the interior of a nuclear cluster.

Qualitatively, the situation was formulated in survey (see Ref. 318): “A nucleus presents not a gas of nucleons, but a pionic soup.” The exchange of such pions at large distances will lead to the interaction between nucleons which strongly differs from the interaction between nucleons in an empty space. These changes lead to both a new expression for the contribution of the strong interaction to the binding energy of a nuclear system and a modification of the Weizsäcker formula which acquires additional parameters. It includes, besides the ordinary parameters A and Z , the fractal dimensionality of the surface D_f , and the parameter of nonextensiveness q . In this case, changes in the binding energy make, naturally, the main contribution in the region of large mass numbers and induce the appearance of new stable states. The Weizsäcker modified formula is analyzed in Sec. 11.1.

12.1.6. Nucleosynthesis in Nature and in a Laboratory: Idea of the Processes of Nuclear Combustion of a Substance

Recently, the Nature gave the materials of the unique experiment supporting the theoretical pattern of nucleosynthesis (see Ref. 319). In 1987, the burst of a supernova in the nearest galaxy “Great Magellanic Cloud”, being at the distance of 180 000 light years, was registered by astronomers. The supernova named CH 1987 A was discovered on February 23, 1987 and underwent a continuous thorough observation with the help of all the available tools. We mention the following registered extremely important processes.

According to the estimates, the collapse continued approximately 2,00 ms. When the falling of a substance to the center stopped, the powerful shock wave was formed for the time interval of 0.4 ms. This wave propagated from the center outward with the initial velocity of about 5×10^9 cm/s.

For the first time, the powerful neutrino emission at the collapse moment was discovered, the existence of the synthesis of heavy elements in the bursts of supernovas was proved, the source of energy of the emission of the supernova shell (the decay of radioactive ^{56}Co) was established, and the electromagnetic emission of the supernova in the whole spectral range (from radiowaves to γ -emission) was studied.

An interesting feature of the supernova burst was the observation of two bursts of the neutrino emission separated by the time interval of about 5 h. According to the commonly accepted scenario, the first burst was related to the generation of a neutron star. The majority of the gravitational energy released in the course of the collapse (10^{53} erg) is spent for the creation of neutrinos in the process of neutronization of the substance. Neutrinos escape outward for the time of $1 \div 10$ s, and their characteristic energies are of the order of 10 MeV. The energy, which is released in the course of the collapse and is expended for the formation of a shock wave breaking away the shell, is lesser by one order.

By the results of measurements, we may conclude that the energy of the second burst has the same order. We can assume that the energy source for the second burst is the phase transition of nuclei of the neutron star in the state of π -condensate. The estimates indicate (see Ref. 300) that the time interval, during which the neutron star is significantly cooled at the expense of the neutrino emission, equals to several hours. As a result of a decrease in the temperature, the equation of state of the nuclear substance approaches the van der Waals one, and a phase transition of the first kind into the superdense state of π -condensate, which occurs for the time of about 10 s, becomes possible. These numbers well agree with experiment.

As a very important phenomenon registered in the laboratory investigations carried out at Electrodynamics Laboratory "Proton-21", we mention the emission arising as a result of the kinetic processes initiated by a REB in a target (see Ref. 209). The idea of a state of the substance in that region of a target, in which nuclear processes are running, is given by the analysis of the emission from the region with extreme parameters (HD) in the X-ray spectral region. The distributions of spectral density of the emission in the X-ray- and γ -ranges from HD of a target (see Ref. 209), those from the astrophysical objects such as a pulsar from the Crab nebula, quasar 3C 273, and short-time splashes of γ -emission are very similar and have the very high level of correlation lying in the scope of $0.92 \dots 0.99$. This can testify

to the existence of common physical mechanisms causing the formation of X-ray emission in astrophysical objects and in the regions with extreme parameters in the experiments performed at the Electrodynamics Laboratory “Proton-21”.

The typical view of the distribution of spectral density of this emission is presented in Chapter 4.

As usual, the spectra include four different sections: two exponential and two power ones. At present, the analysis of possible physical mechanisms of generation of the X-ray emission from a target does not allow us to make the unambiguous choice of some mechanism among many advanced ones.

12.1.7. Conclusions of the Analytic Survey

The choice of the material presented above in this chapter was not accidental. It was defined by the general conception of the possibility to initiate the collective self-consistent nonlinear processes of nuclear combustion in a condensed medium. This conception underlay the approach to both the search for an optimum scheme of experiments and the construction of theoretical ideas. The results of the studies performed in 1972–1982 in the field of optimization of phase trajectories and the analysis of general principles of the evolution of multiconnected dynamical systems (see Refs. 320–321) gradually led one of the authors to the conclusion that the purposeful realization of nuclear reactions and the optimization of structures (the system of connections between elements) of a multidimensional dynamical system have many common points.

The strong nonlinearity of the system studied by us means, in particular, that a change in the external action on the system can qualitatively change a system state and the direction of its evolution. Thus, the external action on the system must be a key factor for the control over the processes of evolutionary synthesis (self-organization) of nuclei.

In Ref. 209 on the basis of a general universal principle, *the principle of regularization of perturbations and dynamical harmonization of systems* (or, in brief, the principle of dynamical harmonization) (see Refs. 320–321), the preliminary analysis of the problems of choice of the external action in order to realize the synthesis of nuclei was carried out. This analysis leads to a conclusion that the success in the solution of the problem of synthesis of nuclei can be reached upon both the self-consistent use of all eigenmodes of motions of the dynamical system and the impact excitation of all its degrees of freedom. Such an action on the system must induce the excitation of all modes of oscillations and all dynamical components of the system with maximally possible amplitudes, i.e., of all derivatives with

respect to coordinates up to the maximum order defined by the dimension of a dynamical system.

By analyzing the conception, we conclude that the nucleosynthesis under experimental conditions is possible as a self-organizing process of pycnonuclear regeneration (combustion-evaporation) of any initial substance.

To realize pycnonuclear reactions, it is necessary to deliver the energy of the order of a Coulomb barrier onto nuclear scales. Since the main factor in the realization of pycnonuclear reactions in the scope of the traditional scenario of their running turns out to be the substance density, we can assert that the pycnonuclear reactions are the nuclear stage of the development of hydrodynamic one-fluid processes.

We note that an important feature of the Kiev experiments consists in that the processes of nucleosynthesis run against the background of a strong current passing along the system. The evolution of magnetohydrodynamic processes related to this current leads to the appearance of plasma-field nonlinear structures (in particular, vortex ones) with extreme values of the magnetic field, which significantly affects the nuclear processes.

As a natural generalization of one-fluid processes, we indicate two-fluid flows of a substance with electrodynamic effects intrinsic to them.

Upon the realization of these modes, the nuclear stage runs as a result of the realization of the electrodynamic version of pycnonuclear reactions. In the last case, the self-consistent electromagnetic field unseparably connected with the two-fluid scheme plays the important role, as well as with the substance density. Such an establishment of the conditions for the running of pycnonuclear reactions can significantly enhance their efficiency. In electrodynamic pycnonuclear reactions, the essential role should be played by the strongly nonlinear processes of polarization of the medium under significantly nonequilibrium conditions.

The two-fluid plasma stage is accompanied by the development of plasma-field structures. For their excitation, we need, in particular, to create a nonlinear wave of volume charge which attains the extreme parameters in the course of its evolution and can initiate a plasma-field structure. The last has high density and involves self-consistent electromagnetic fields on the atomic and nuclear scales as for their values and the degree of their localization.

The wave of volume charge is most efficiently excited by beams of charged particles which serve simultaneously as the efficient sources of ordinary one-fluid hydrodynamic flows.

In the experiments performed at the Electrodynamics Laboratory "Proton-21", the primary driver of the compression of a target was a heavy-current relativistic electron beam. For the superstrong compression of

a target, it is necessary, as known, to realize the sharpening modes. Prior to the Kiev experiments, the sharpening modes were realized only in the laser thermonuclear synthesis. Experiments on laser inertial synthesis were oriented to the realization of the one-fluid hydrodynamic scheme. On the contrary, the experiments performed at the Electrodynamics Laboratory "Proton-21" have included, for the first time, the sharpening mode on a REB, as well as the two-fluid mode. To realize the excitation of a wave of volume charge in the high-conduction medium is not a simple task. In our experiments, it was solved due to the realization of a specific current transfer along the system.

In the following section, we will construct a qualitative model of electrodynamic processes in a target and give the description of the anode as an element of the electrical delay line rather than a localized element of the setup. In the case where a current pulse is of poor quality (e.g., the front duration is not sufficiently small), the delay line behaves itself as the element with lumped parameters, and the charge accumulation effect at the beginning of the line is lacking.

In view of the fact that a pulse action with small duration has maximally wide spectrum, namely the pulse action is suitable for the efficient initiation of evolutionary processes. The type of a pulse action should correspond to the principle of excitation of a maximally wide class of types of the natural oscillations of the medium. An external pulse action have to initiate all the internal resources of the system.

As mentioned above, according to the results of theoretical investigations (see Ref. 233), the centrally symmetric compression of a finite mass of a substance up to extreme densities can be executed with the use of a spherical piston moving with a certain acceleration. This conclusion agrees with that following from the evolution-based principles.

Prior to the Kiev experiments, a "piston" for the inertial synthesis was formed by a special inhomogeneity of a target. A distinct feature of the Kiev experiments is the use of self-consistent hydrodynamic (and electrodynamic) modes for the creation of a piston moving to the center, namely a thin quasispherical layer with high density (a wave-shell).

Within this approach, we may hope for the success even with low initial energy resources. It is only necessary to concentrate the energy of a driver on the particles of a substance, whose number should be coordinated with this energy.

Further, we present the theoretical substantiations of the above-discussed general positions of the conception concerning the realization of the self-organized nucleosynthesis.

In Sec. 12.2, we discuss the main ideas of peculiarities of the formation of the electron beam in a relativistic diode constructed at our laboratory. We give the scenarios (the one-fluid and two-fluid modes) of the appearance of extreme parameters in a target under the action of a pulse REB and construct theoretical models for the development of nonlinear processes in a target which induce plasma-field structures (in particular, vortex ones) in the one-fluid and two-fluid modes.

In Sec. 12.3, we study peculiarities of the formation of dynamical nuclear structures with a self-consistent account of the disequilibrium of the system. We show the possibility to form nuclear clusters and analyze the kinetics of their formation and peculiarities of the dynamics of clusters. We analyze the plasma-field structures in nuclear clusters and on their periphery in external self-consistent fields under the electron and pion condensations. A Weizsäcker modified formula is deduced, and the binding energy of clusters and their dynamical stability and evaporation are analyzed.

In Sec. 12.4, we summarize the studies performed and outline the perspectives of theoretical ideas and possible experimental realizations of the electrodynamic pycnonuclear reactions and the processes of self-organized nucleosynthesis. In particular, we analyze all possible developments of some traditional approaches to the CTF from the viewpoint of our conception.

12.2. The Theory of Energy Concentration on Nuclear Scales

As shown in Sec. 12.1., a key problem in the initiation of collective nuclear processes in the condensed medium is the concentration of the energy of a primary driver on an optimum amount of a substance on a sufficiently small spatial scale.

Below, we formulate the requirements that should be satisfied for the efficient solution of the energy concentration problem on a small scale:

- The delivery of energy on nuclear scales must be realized by a nonlinear process,
- This process should be accompanied by a decrease in the Coulomb barrier of nuclei,
- In view of the electromagnetic nature of the repulsive forces, which must be overcome, the nonlinear process should be related to charged components of the medium, namely electrons and ions.

Pulse relativistic electron beams (REBs) used in the Kiev experiments correspond to all these requirements. As their characteristic feature, we indicate that a REB can simultaneously efficiently perturb all the components of a substance: the neutral component, electrons, ions, and natural electromagnetic oscillations of the medium.

In the frame of the traditional schemes of inertial thermonuclear synthesis, the attempts to use a REB as a driver were stopped because of the problems of focusing. Though a high degree of the focusing of REBs was achieved in the best experiments some time ago, but this focusing was unstable. Moreover, a symmetric squeezing of a target was not realized (see Refs. 217, 231) on the stage of the one-fluid dynamics of a quasineutral medium in the bulk of a target.

Due to the intense experimental and theoretical studies of the processes in hard-current diodes performed at the Electrodynamics Laboratory “Proton-21” (see Ref. 209), we succeeded to solve this complex problem and to reach a symmetric squeezing of the required part of the anode.

From our viewpoint, the internal degrees of freedom of complex multidimensional nonlinear dynamical systems with interactions, which cannot be regarded as small perturbations, were not used to a full extent in the inertial synthesis (despite a considerably higher level of the self-consistency of processes as compared to that with applying the methods of the force keeping of a plasma in the thermonuclear synthesis).

The block diagram in Fig. 12.19 displays the interconnections of main processes running in the anode upon the action of a REB.

For the description and comprehension of the essence of processes running upon a pulse action of electron beams on a solid surface, we developed the conception of the interconnection of plasma-beam and nuclear processes. The first version of this conception is given below.

As was expected, the main difficulty in the development of the conception consisted in the necessity of a self-consistent description of collective processes in many-particle systems. We now consider the macroscopic models for all main dynamic processes shown in Fig. 12.19, beginning from models of a relativistic diode.

12.2.1. Model of a Relativistic Diode with Plasma Electrodes

The problem of formation of the hard-current beams of charged particles is associated with the solution of a number of complex self-consistent nonlinear electrodynamic problems. Computational difficulties are defined by strong nonlinear effects, instable modes in a diode, the possibility of a partial “locking” of beams by the field of the own volume charge, and the arising of the mode of a virtual cathode under the availability of large currents in a diode.

The analysis of pulse hard-current setups of various types allows us to separate the main elements and radiophysical parameters necessary for their simulation. Main elements include a capacity storage unit, a storage unit with lumped and distributed inductances, a plasma breaker of current, a plasma-filled hard-current diode as a load, the totality of elements of a circuit

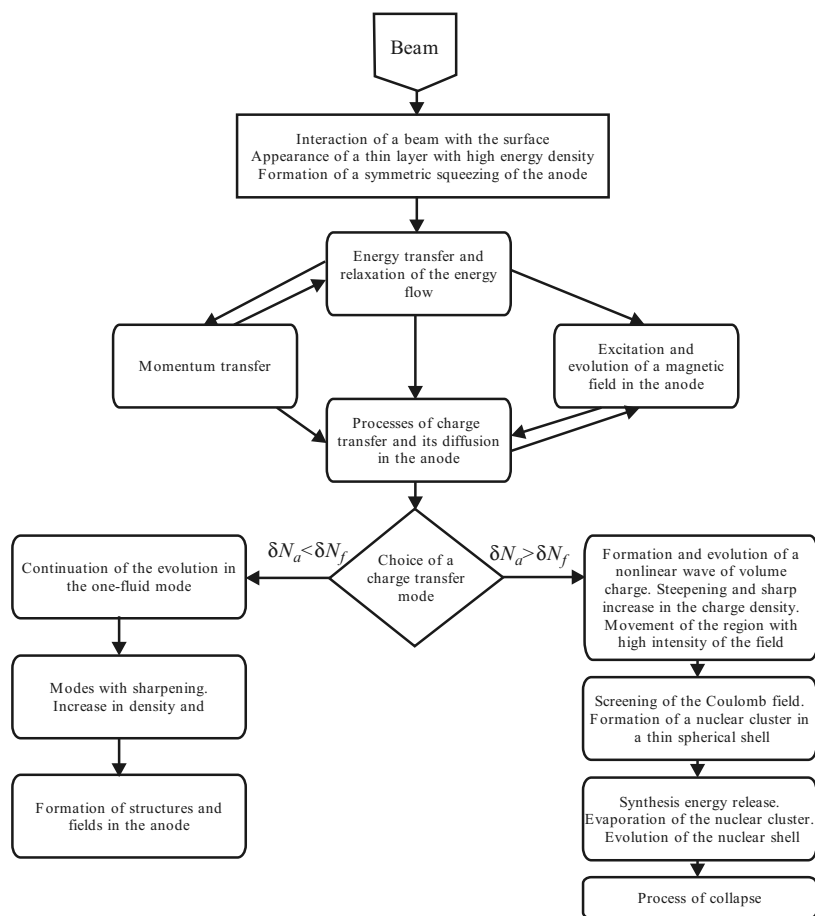


Fig. 12.19. Block diagram of the interconnection of main processes upon the laboratory nucleosynthesis. is the excess of the density of electrons in the anode as a result of the action of a REB. is the amplitude of fluctuations of the density of electrons.

for the account and calculation of various losses and the inverse current, i.e., the current from the diode anode to the ground electrode of the storage unit (see Refs. 236, 322–323).

The typical equivalent scheme of a relativistic diode, which has all above-indicated components and was developed by the Micro-Cap7 system of scheme-technical simulation (see Ref. 324), is shown in Fig. 12.20.

This scheme demonstrating a general structure of a hard-current diode should be filled by a physical content qualitatively corresponding to the main processes running after the creation of conditions for the initiation of plasma-involved and nuclear processes in the anode bulk.

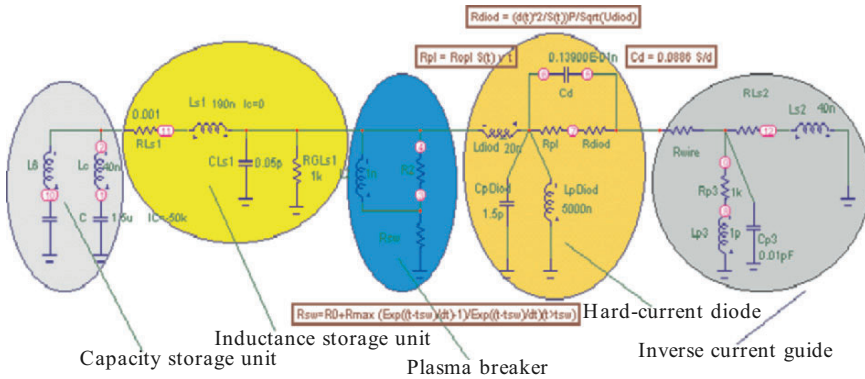


Fig. 12.20. Block scheme of a hard-current diode.

In the modern pulse technique, the main method to get hard currents in a diode is based on the use of cathodes with explosive emission of electrons. The explosive electron emission (EEE) consists in the yield of electrons from a metal in the process of explosive phase transition of a cathode substance from the condensed state to a dense plasma. Such a phase transition can be realized under a fast heating of a local section of the cathode up to a high temperature, at which its explosion-like evaporation and ionization begin.

This situation is realized where, due to the supply of a high voltage on the vacuum gap, there occurs the explosion of microscopic points on the cathode under the action of the passing thermo- and autoemission current (see Ref. 325). Local plasma clusters formed on the cathode are propagating in vacuum. Conduction electrons pass from a metal through the zone of the phase transition “metal–dense plasma” in the cathodic flare (CF) and are emitted from the moving plasma boundary in vacuum. The transition from thermo- and autoemission of electrons to the mode of explosive emission is accompanied by increasing the emission current by almost two orders.

The plasma velocity in small-size gaps weakly depends on the applied voltage, is practically constant in time, and equals $(1 \div 3) \times 10^6$ cm/s for all metals. The principal elementary process is the ionization of atoms by an electron impact. The temperature of electrons in the plasma of a cathodic flare is $4 \div 5$ eV, and the concentration of electrons at the distance of 10^{-2} cm from the cathode surface is of the order of 10^{16} cm $^{-3}$.

The dispersion of a plasma is satisfactorily described by the hydrodynamical model of expanding CF. Suppose that the size of a flare becomes much more than the initial one of the exploded metal. By assuming that the plasma dispersion conditions are close to adiabatic, we can get the plasma propagation velocity in the flare as

$$v = \sqrt{\frac{4\gamma}{\gamma - 1}(\varepsilon_0 + Z\varepsilon_F)}, \quad (12.167)$$

where γ is the isentropic exponent; ε_0 is the initial specific energy equal to the sublimation energy of a cathode material; Z is the charge number of an ion; ε_F is the Fermi energy. The value of v calculated by this formula well agrees with the experimental one. This model is corrected with regard to the ionization and heating of electrons of a plasma CF at the expense of the Joule mechanism of energy dissipation upon the passage of current through the plasma. It was shown experimentally that the region of the efficient release of energy is near the cathode, and there occurs the energy transfer from internal regions of the plasma to the periphery.

For various geometric situations, it was also demonstrated that the law of three halves holds in the diode with a moving plasma cathode (see Ref. 325).

The specificity of the setup is defined, first of all, by the parameters of a breaker and a diode. Starting from the first principles, we have developed the models of the operation of plasma breakers and hard-current plasma diodes describing their work and the dependences of their parameters on time (see Refs. 326, 327). The parameters of breakers are the pickup times t_{sw} , characteristic times of a variation of the resistance τ_{sw} , initial and final resistances R_0 and $R_{\max}$, t_r is the time moment of the beginning of the decay of a plasma-field structure, τ_r is the characteristic decay duration of the structure. The law of variation of the breaker resistance can be set analytically. One of the admissible models of breaking is the exponential law of variation of the resistance, for example,

$$R(t) = R_0 + (R_{\max} - R_0) \left(\frac{\exp(\frac{t-t_{sw}}{\tau_{sw}}) - 1}{\exp(\frac{t-t_{sw}}{\tau_{sw}})} \right)^2 \times \left(\frac{\vartheta(t - t_{sw})}{1 + \exp(\frac{t-t_r}{\tau_r})} \right). \quad (12.168)$$

This function of time is displayed in Fig. 12.21.

A peculiarity of the diode is the appearance of a plasma cathode in it.

The parameters of the plasma cathode are the following: the speed of motion of a plasma v_{pl} , the area of the emitting surface depending on time $S_{pl}(t)$, a moment of the appearance of a plasma in the diode t_{pl} (it is physically related to the peculiarities of the discharge over electrodes in the diode), the plasma conduction σ_{pl} (defined by the density n_{pl} and temperature T_{pl} of the plasma). To adequately describe the operation of the diode, it is necessary to take into account both the dynamics of the plasma



Fig. 12.21. Resistance of a plasma breaker versus time.

cathode and the movement of the anodic plasma. Elements of the dynamics of this plasma are considered in the next section.

A minimum collection of the parameters describing the plasma cathode includes a time moment of the beginning of the motion t_{pl} and the velocity v_{pl} of the cathodic plasma. Based on the geometry of the diode and on physical models, we can define the time dependences of the area of the emitting surface of the plasma cathode $S_{pl}(t)$ and the varying gap between the plasma cathode and the plasma anode $d(t)$.

We note that the time moment of the start of motion of a plasma in the diode can differ from the time moment of a breaking of the current by a plasma breaker depending on the construction of a setup. These differences are among the important parameters of the setup. The proper choice of the time moment of a breaking by a plasma breaker is very important for the efficient functioning of the setup. In this case, the resistance of the anode - cathode gap drops by the law

$$R_{diod}(t) = \frac{d(t)^2}{S_{pl}(t)P} \frac{1}{\sqrt{U_{diod}(t)}}, \quad (12.169)$$

where P is a characteristic value of the diode perveance.

A decrease in the resistance of the anode-cathode gap is accompanied by a growth of the additional resistance of the plasma bridge $R_{pl}(t)$. This resistance changes from zero to a small value defined by the resistance of the initial clearance d_0 completely filled by a plasma.

The minimum capacitance of the diode can be estimated by the formulas for a planar or coaxial capacitor; a real capacitance is defined by structural features of the diode. In Fig. 12.22, we show an example of the characteristic dependence of the main parameters of the setup on time.

We now pass to a brief theoretical analysis of the main plasma-beam processes in a hard-current diode.

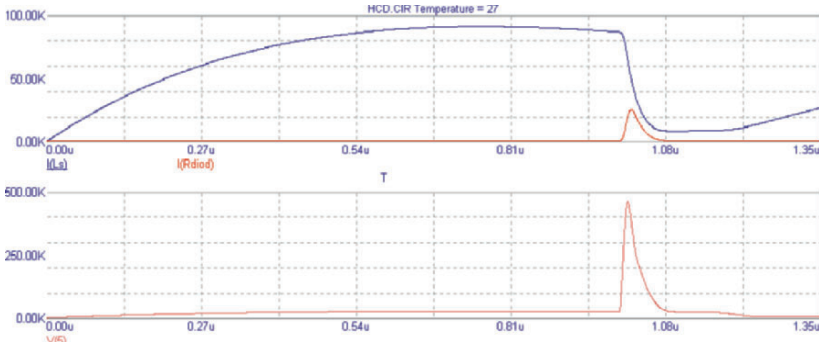


Fig. 12.22. Main results of calculations of the equivalent scheme of a diode. In the upper figure, we show pulses of the current of an inductance storage unit and the current in the diode. The lower curve presents a voltage pulse appeared as a result of the breaking of the current by a plasma breaker.

12.2.2. A Hydrodynamic Theory of Electron Beams and an Anodic Plasma in a Diode

The analysis of the dynamics of beams and a plasma in a hard-current relativistic diode is necessary, because just a conscious use of some nonlinear and nonstationary phenomena in the diode allows one to initiate the self-consistent processes leading to the concentration of energy in a target, which is necessary for the realization of nuclear processes.

An analytic model of the fluxes of charged particles (see Ref. 328) was developed with the use of small parameters related to the geometric features of diodes (thin electrodes and diodes with a large aspect ration).

The basis of the theoretical model describing the dynamics of hard-current relativistic beams in cylindrically symmetric electrodynamic systems is the hydrodynamic equations for an electron flux in a cylindrical coordinate system (r, θ, z) .

This system of equations includes the Euler equations for an electron flux [these equations take also into account external electric and magnetic fields governing the focusing $E_z^{ext}(r, z)$, $E_r^{ext}(r, z)$, $H_\theta^{ext}(r, z)$], equations of continuity for the electron fluid density, and system of the Maxwell equations for vortex $(\vec{E}, \vec{H})$ and potential fields $\vec{E}_p$.

To get an analytic solution of the system under study, we use the method of perturbation in a small parameter $\eta = \frac{R}{l} \ll 1$, where R is the characteristic scale of variations in the transverse coordinate and l is the characteristic longitudinal size of the system in the case of needle electrodes or in the inverse parameter $\eta = l/R \ll 1$ for diodes with large aspect ratio.

We will seek the solutions of the system of equations by using the representation of all functions F and other quantities in the form of expansions in a small parameter η :

$$F = F^0 + F^1\eta + F^2\eta^2/2 + F^3\eta^3/3! + \dots \quad (12.170)$$

Collecting the coefficients of the identical degrees of η in the relevant equations, we get the systems of equations for the successive approximations to solutions $F^0, F^1, F^2 \dots$.

We note that the problem is multiparametric, and dimensionless parameters appear upon making all the physical quantities to be dimensionless with regard to their physical dimensionality and the type of a source of energy defining the dynamical processes in the system.

In particular, by solving our problems, we need to distinguish electric fields of different nature (an external field, potential radial and longitudinal fields, vortex fields, etc.) and to make them dimensionless with the use of different quantities according to their essence. For example, we make an external longitudinal field to be dimensionless by dividing it by its maximum value, and the internal potential fields become dimensionless with the use of the initial values of the energy and density of a beam.

The dimensionless system of equations consists of the equations for components of the vortex electric field of a beam, magnetic field of a beam, and potential electric field of a beam, equations of continuity, and equations of motion. The derived equations with regard to the expansion of all quantities in series Eq. 12.170 with the use of algorithmic procedures written in the Mathematics 5 language yield the equations of zero, first, and higher orders in a small parameter (see Ref. 328).

We note that the dimensionality of these systems increases rapidly with the order of exactness in a small parameter. Their use is justified only in the automated media of analytic transformations, and therefore we do not give their specific form there.

The analysis of the systems of zero and first orders allows us to clarify the successive appearance of various spatial components of the fields in the electrodynamic system and their influence on the dynamics of particles.

In particular, under conditions of the explosive emission from the cathode, the distribution of the current over a radius is approximately described by elliptic functions, and two spatial components of the electron beam are clearly observed: the component propagating along the system axis and the component having the form of a cylindrical shell on the periphery of the system (in the simple case of axial symmetry). As a result, the analytic theory of diodes with plasma cathodes of arbitrary forms is actually constructed.



Fig. 12.23. Structure of the electron beam upon the explosive emission from the cathode.

The derived two-component structure of an electron beam in the diode with a plasma cathode is of importance for the excitation of subsequent processes on the anode and is related to the phenomena of magnetic self-isolation of a beam. The two-component spatial structure of electron fluxes in a diode with explosive emission was confirmed experimentally (see Fig. 12.23).

It is known (see Refs. 236) that plasma layers arise upon the generation of a REB in hard-current diodes on its electrodes and on the walls of a chamber. The movement of a cathodic plasma has already discussed above. We note that a lesser attention is usually paid to anodic plasma, as compared to cathodic one.

In our case, the initiation of nonlinear waves in the anode with high symmetry (cylindrical or spherical) is defined by features of the motion of both the cathodic and anodic plasmas.

The layers of anodic and cathodic plasmas in a hard-current diode serve, respectively, as sources of positive ions and electrons. In a diode, electrons form a REB, and ions form the fluxes directed toward the cathode. These ion fluxes render a positive effect on the running of processes in diode-based devices, in particular, on an increase of both the emission of electrons and the total limiting current of a diode.

For hard-current diodes, the question on the dynamics of ions is also important from the viewpoint of the determination of relations between the ion and electron components of a current. The presence of fast ion fluxes establishes the conditions for the creation of a beam and the ion medium providing the compensation of a volume charge, which is necessary for the focusing of a REB, on the earlier stages of the evolution.

The anodic plasma, whose source is the anode bombarded by the electron beam, emits ions. Upon the action of electric fields, ions propagate

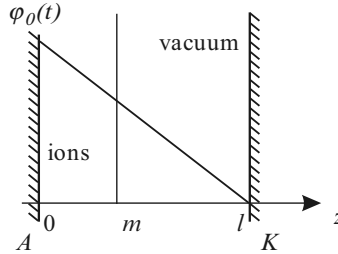


Fig. 12.24. Electrodynamic structure of a cathode–anode gap.

in the region near the anode, including the direction towards the cathode. In addition, it turns out that the own volume charge of the ion cloud renders a significant effect on the increase of the velocity of ions (see Refs. 329, 330). The influence of collective effects on the velocity of ions is essential and should be considered, in particular, from the viewpoint of the focusing of REBs. The classical investigations of the dynamics of electron fluxes (see Ref. 331) carried out on the first stages of the development of the physics of diodes are applied to smooth variations in voltage (in fact, to a quasistationary external pulse). The analysis of the dynamics of a current in a diode performed in Refs. 329, 330 allows one to analyze features of the physical processes in a diode depending on the voltage front steepness.

Following the above-mentioned works, we consider a layer of ions (Fig. 12.24) in the Cartesian coordinate system OXZ which arises on the anode upon the supply of a voltage with a sharp front $\varphi_0(t)$ to the cathode–anode gap and propagates from the anode to the cathode upon the action of external electric fields and electric fields of the own volume charge.

Beginning from the time $t = 0$, the planes $z = 0$ and $z = l$ are at the potentials $\varphi(0) = \varphi_0(t) > 0$ and $\varphi(l) = 0$, respectively, so that the plasma is acted by an external electric field at $t > 0$. Then, by the time t close to $t = 0$, ions have shifted from the initial position, and all the region $(0, l)$ turns out to be divided into two physical regions:

- (a) $(0, m)$ – ions,
- (b) (m, l) – vacuum.

We will consider the electrodynamic problems for each region with regard to a nonlinear motion of ions, sew together fields and potentials, and satisfy the boundary conditions.

The hydrodynamic description of ion fluxes is based on the equations of motion of ions,

$$\frac{dv}{dt} = \frac{Ze}{M}E, \quad \frac{\partial z}{\partial \tau} = C_1(\tau), \quad (12.171)$$

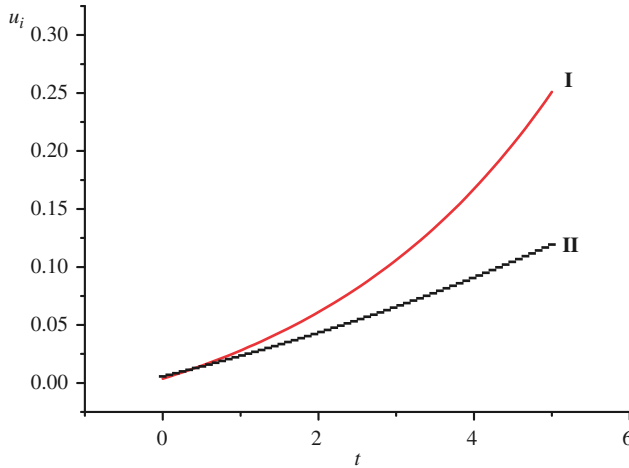


Fig. 12.25. Curve I is the velocity of the ion flux front; curve II is the velocity of a single ion.

the equation of continuity, and the Poisson equation for a self-consistent electric field

$$\frac{\partial E}{\partial z} = 4\pi Z e n, \quad E = -\frac{\partial \varphi}{\partial z}, \quad (12.172)$$

where v, n, E , and φ are the velocity, density, electric field, and potential, respectively.

In system Eqs. 12.171–12.172, t and τ are the Lagrange variables, $C_1(\tau)$ is the function derived after the integration of the equation of continuity described in the Lagrange variables over time. The integration of this system allows us to calculate the velocity of the ion flux front (curve I in Fig. 12.25) and its coordinates (Fig. 12.26).

Figure 12.25 shows the excess of the velocity of the ion flux front above the velocity of an appropriate single ion (curve II). This excess testifies to the presence of the effect of collective acceleration of ions for the supplied voltage pulses with large amplitude and sharp front.

The magnitude of the effect of self-acceleration of ions can be optimized. This effect depends significantly of the growth rate of a pulse voltage on the diode. The effect of self-acceleration increases with the growth rate of a pulse. However, if the rate is very large, it is difficult to get a propagating beam with sharp front. This leads to the existence of the optimum.

It is seen that the velocity acquired by the flux of ions at the expense of collective processes exceeds significantly the velocity of a single ion (i.e., an ion accelerated by only the anode–cathode voltage). The processes of collective acceleration lead to that ions fill rapidly the anode–cathode region (Fig. 12.26).

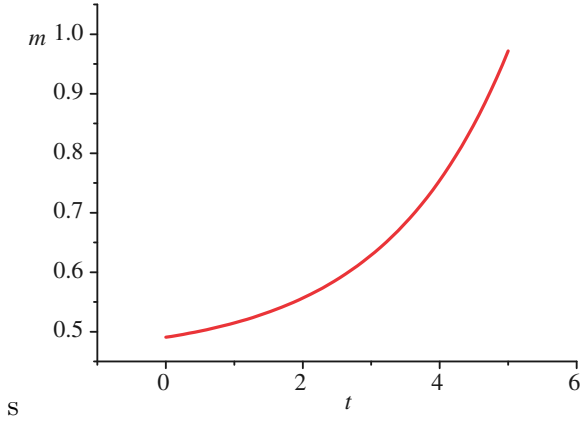


Fig. 12.26. Position of the leading edge of an ion flux.

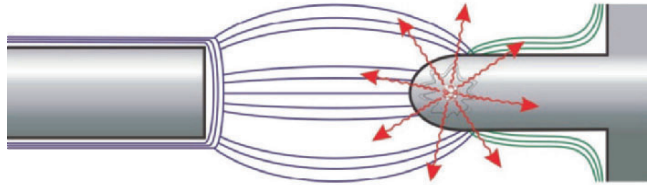


Fig. 12.27. Scheme of the arrangement and the dynamics of the cathodic plasma, anodic plasma, and electron beam.

Thus, there occur the self-acceleration of the ion flux in the diode under the action of a high-voltage pulse and its fast propagation towards the electron beam being focused. An increase in velocities of ions influences positively the focusing of the electron beam at the stages earlier as compared to those in the case of ions considered in the one-particle approximation.

The positions of the cathodic plasma, anodic plasma, and electron beam are schematically shown in Fig. 12.27.

In Fig. 12.28, we displayed the typical form of the pulses of current and voltage derived as a result of the analysis of the equivalent scheme of a diode with its characteristic parameters and with regard to the motion of the cathodic and anodic plasma in the diode.

Upon the focusing of an electron beam, the collective effects analyzed above lead to the appearance of a potential well for ions near the anode, to which the ions are captured, by forming a dense bunch. The density of ions in the bunch is of the order of the density of focused electrons of a beam.

With increase in the current, we observe the focusing of both the spatial components of the electron beam and their approach to the axis. Even for a cylindrical anode, these peculiarities lead to both the appearance of the effective plasma surface, which has the form of a cone with small

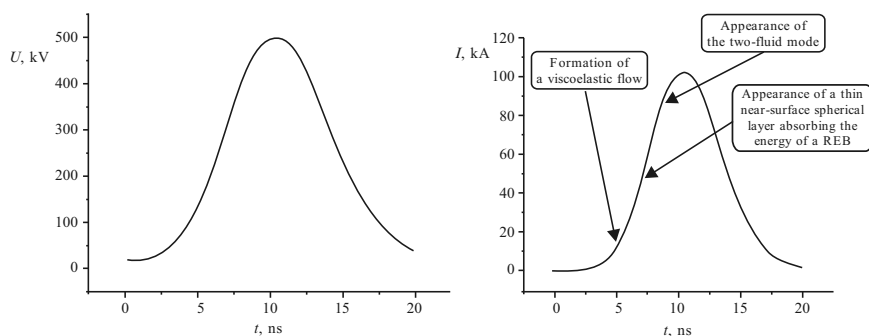


Fig. 12.28. Typical pulses of voltage and current in a hard-current diode.

angle at the tip (see Fig. 12.27), and the fast movement of the region undergone the action of its external cylindrical component towards the cathode upon the scanning of the anode surface by the electron beam.

Such a geometry of the squeezing of a target due to the presence of the longitudinal and transverse components of pressure induces the cumulative effect. This effect significantly increases the pressure on the lateral surface (see, e.g., Ref. 373) and promotes the symmetrization on an external action on the target.

Recoil ions and the electron beam entering the anode and the anodic plasma induce a pressure on the anode. This time-dependent pressure is a source of energy, first of all, for the electron component of the medium. The estimate of this pressure will be given below.

12.2.3. Characteristic Features of the Operation of a Relativistic Pulse Diode with Plasma Electrodes and the Excitation of Nonlinear Waves in a Condensed Medium in the One-Fluid Approximation

Here, we will show how the above-considered features of the nonlinear dynamics of an electron beam and a plasma in a relativistic diode lead to peculiarities of the formation of modes with sharpening in a target and to the excitation of nonlinear waves with extreme parameters.

In order to analyze the processes in the anode under the action of concentrated fluxes of energy, it is necessary to know the equation of state of a substance. As idealized equations of state of a substance, the majority of studies used those for a plasma or the ideal gas. The numerous numerical methods allow one to realize, with various degrees of exactness, the complete calculation of shock waves, contact discontinuities, and their interaction between themselves, and nonlinear waves with the use of “simple” equations

of state of a substance. However, simple models of the equation of state of a substance cannot be used under extreme conditions. In this case, the efficient numerical simulation can be realized only with the use of a wide-range equation of state of a substance which is valid for both a rarefied plasma and for a solid experienced an extreme pressure.

A Wide-Range Interpolational Equation of State of a Substance.

Below, we present an interpolational equation of state proposed for copper in Ref. 333. According to this equation of state, the total pressure $P(n, T)$ and energy of a substance E consist of several components and are defined as follows.

The total pressure is defined by the relation:

$$P(n, T) = P_e + P_0 + P_{el}, \quad (12.173)$$

where the pressure of electrons

$$P_e = 0.667n_e E_e, \quad (12.174)$$

pressure of ions and neutral atoms

$$P_0 = \frac{\gamma + a/3}{1 + a/2} E_0 n_s, \quad (12.175)$$

and the elastic component of the pressure

$$P_{el} = 0.00195n_s^3 - 0.0376n_s^{7/3}. \quad (12.176)$$

The specific energy per unit volume of a substance E is also defined by the sum of several terms as

$$E(n, T) = E_e + E_0 + E_{el} + E_i. \quad (12.177)$$

In these relations, E_e , E_0 , E_{el} , E_i are, respectively, the energy of electrons, energy of ions and neutral atoms, elastic strain energy, and energy spent on the ionization of a substance. These quantities are defined by the relations

$$\begin{aligned} E_e &= 0.457\varepsilon_m Z \ln(ch(3.27T_e/Z\varepsilon_m)), \\ E_0 &= \frac{2+a}{1+a} 1.5 T_0, \\ E_{el} &= 2.12 \left[2 \left(\frac{n_s}{85} \right)^2 - 3 \left(\frac{n_s}{85} \right)^{4/3} \right], \end{aligned} \quad (12.178)$$

where $a = 673.46T_0/n_s^{1.479}$, $\gamma = 1.917$, and $\varepsilon_m = 0.364n_e^{2/3}$.

Here, we use the following designations: ε_m is the maximum energy of a degenerated electron gas in a crystal (see Ref. 261), n_j is the concentration of ions of multiplicity j , $n_s = \sum_{j=0}^L n_j$ is the total concentration of heavy particles (ions and atoms), L is the nucleus charge, I_j is the j th ionization potential, Z_s is the equilibrium ionization degree by Saha, and γ is the Grüneisen coefficient. In these relations and the equation of state, the energy is given per one particle, the concentration is measured in units of 10^{21} cm^{-3} , pressure in units of 1.602 kbar, and temperature and energy are given in eV.

The energy spent for the ionization of a substance,

$$E_i = \frac{1}{n_s} \sum_{j=1}^L \left(\sum_{k=j}^L n_k \right) I_j, \quad (12.179)$$

is naturally defined by the ionization potentials I_j which depend on the ionization potentials for a rarefied plasma I_{0j} and the density of a substance as

$$I_j = I_{0j} \left\{ 1 - \ln \left[1 + \exp \left(4.55 - 6.37 \sqrt{I_{0j}/n_s^{1/3}} \right) \right] \right\}. \quad (12.180)$$

In Fig. 12.29, we present the ionization potentials for metallic copper under normal conditions, and Fig. 12.30 shows the multiple-ionization potentials of copper versus density. As seen, the ionization potentials of all multiplicities sharply drop with increase in the substance density. Such

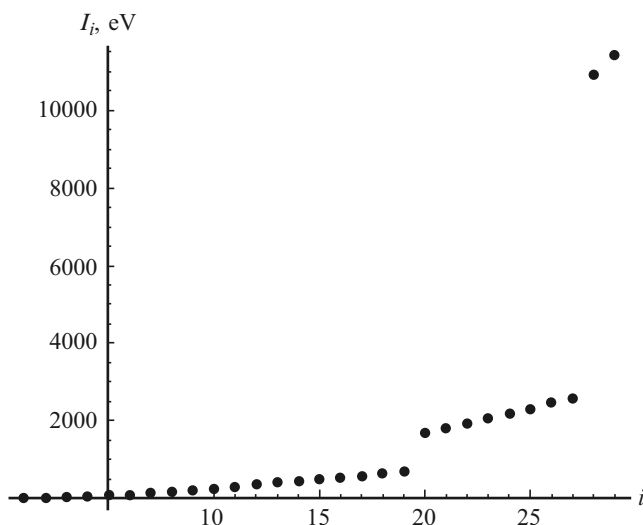


Fig. 12.29. Ionization potential of copper (eV) versus the ionization degree from 1 to 29.

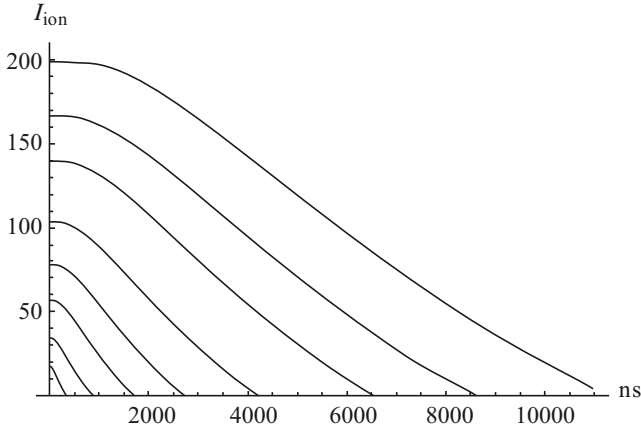


Fig. 12.30. Dependence of the first ionization potentials of copper on the excess of the density above the normal value.

a decrease in the ionization potential should be necessarily taken into account in the analyzed processes running in a target, because the density is locally increased in a wave propagating in a target by several orders.

The ionization degree of a substance Z and the density of multiply charged ions n_j are defined from the relations

$$Z = \sum_{j=1}^L j n_j / n_s, \quad (12.181)$$

$$n_j = \left(\frac{n_{j-1}}{n_e} \right) K_j(n_s, T_e), \quad (12.182)$$

$$K_j(n_s, T_e) = 6 T_e^{1.5} A'_j \exp(-I_{0j}/T_e), \quad (12.183)$$

where $A'_j = 1 + \frac{10n_s}{T_e^{1.5} + (10I_j)^2 \exp(I_j)}$. The correction A'_j characterizes a deviation from the equilibrium by Saha (the higher the density of a substance, the greater the deviation) and provides the convergence of the ionization degree to the Saha formula in the plasma-state limit and to 1 in the solid-state limit.

The dependence of the coefficient $K_j(n_s, T_e)$, which defines the recurrence relation between the densities of ions with the ionization degrees differing by 1, on the density and temperature is given in Fig. 12.31.

It is seen that a high degree of ionization in a target can be attained not only due to the increase in temperature, but also at a low temperature at the expense of a high density of a substance.

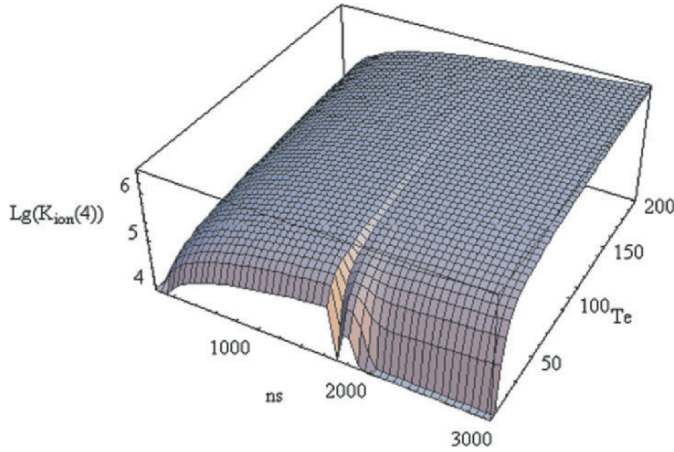


Fig. 12.31. Dependence of the ionization degree of a dense nonideal plasma on the density and temperature (the concentration n_s is measured in units of 10^{21} cm^{-3} , and the temperature of electrons T_e in eV).

Besides the density of ions, one of the main values characterizing the processes in a plasma is the characteristic time of electron-ion relaxation τ_{ei} :

$$\frac{1}{\tau_{ei}^{\sigma, \chi}} = 2 \cdot 10^{12} \left(1 - \frac{B^{\sigma, \chi}}{T_e + \varepsilon_m} \right) \frac{S_0}{A^{\sigma, \chi}} \sqrt{1.5 T_e + \varepsilon_m} \times \left(n_0 + 208 \sum_{i=1}^L n_i \Lambda_i \left(\frac{i}{1.5 T_e + \varepsilon_m} \right)^2 \right). \quad (12.184)$$

Here, $S_0 = 0.02 T_i / (T_{pl} + 0.0027 T_i)$, Λ_i is the Coulomb logarithm, $A^\sigma = 0.65$, $B^\sigma = -3.0$, $A^\chi = 1.74$, and $A^\chi = -9.64$.

The coefficients of electroconductivity and thermal conductivity are simply expressed through the time of electron-ion relaxation τ_{ei} as follows: $\sigma = n_e e^2 \tau_{ei}^\sigma / m_e$ and $\chi = n_e k_B^2 T_e \tau_{ei}^\chi / m_e$.

The formulas for $P(n, T)$ and $E(n, T)$ have such a form that they approach asymptotically the well-known relations of solid-state physics (in the limit of high densities and small temperatures) and those of plasma physics (in the inverse case).

Completing the analysis of the equation of state, we present the formula for the sound velocity. In a medium, perturbations propagate with the adiabatic sound velocity that is defined as $c = \sqrt{(\partial P / \partial \rho)}|_S = \text{const}$.

To analytically determine the sound velocity, we need to know the equation of adiabat, $dS = 0$. In the general case, it has the form $dE + PdV = 0$ or, in our case,

$$d(E_0 + E_e + E_{el} + E_i) = -(P_e + P_i + P_{el}) dV. \quad (12.185)$$

Substituting the well-known expressions for P and E in the last relation, we get the equation of adiabat as

$$\begin{aligned} \frac{5.75 + a}{1 + a} T_0 \frac{dn_s}{n_s} = 1.5 \left[\left(1 + \frac{1}{(1 + a)^2} \right) dT_0 + \frac{1.479a}{(1 + a)^2} \frac{T_0}{n_s} dn_s \right. \\ \left. + \frac{1}{n_s} \sum_{j=1}^L I_j \left(\sum_{i=j}^L dn_i \right) + 1.5 T_e Z \operatorname{th} \left(\frac{3.27 T_e}{Z \varepsilon_m} \right) \right. \\ \left. \times \left(\frac{dT_e}{T_e} - \frac{dZ}{Z} - \frac{2}{3} \frac{dn_e}{n_e} \right) + \frac{8}{3} E_e dZ \right]. \quad (12.186) \end{aligned}$$

For small perturbations and high densities in the adiabatic approximation, the sound velocity (cm/s) can be approximately written as

$$\begin{aligned} c = 1.58 \cdot 10^5 \left[T + 8.09 n^{1/3} \left(\frac{dT}{dn} - \frac{2}{3} \frac{T}{n} \right) + 0.6n \frac{5.75 + a}{1 + a} \left(\frac{dT}{dn} + \frac{T}{n} \right) \right. \\ \left. - 9.57n \frac{a}{(1 + a)^2} \left(\frac{dT}{dn} - 1.479 \frac{T}{n} \right) \right. \\ \left. + 0.6n^{4/3} \left(0.00585n^{2/3} - 0.0877 \right) \right]^{1/2}, \quad (12.187) \end{aligned}$$

where $\frac{dT}{dn} = \frac{5.75}{2} \left(\frac{2}{3} \frac{T}{n} \right)$.

Upon the action of a beam, the density and temperature of a metallic target increase. The substance is ionized, and the ionization equilibrium is defined by the attained temperature and the decrease in the ionization potential as a result of the increase in the target density.

The equations of state of a substance in the form of Eqs. 12.173–12.187 allow us to describe the evolution of the main physical quantities during the action of a beam. Under the action of a beam, a dense plasma arises on the surface of the anode, and the transient layer between the anodic plasma and a metal possesses particular properties (see Ref. 333). Responsible for these unusual properties are the collective interactions of an electron beam with the anodic plasma and the plasma in a solid. These properties manifest themselves especially strongly for the fluxes of charged particles with sharp fronts and great currents. In the next section, we will consider the main physical effects appearing under the collective interaction of a beam of electrons with a target and a macroscopic model of the nonstationarity of the impedance.

Collective Processes of Loss of the Beam Energy in a Target and a Nonlinear Model of the Evolution of the Impedance of a Relativistic Pulse Diode. We describe the main processes which lead to the formation of a region of losses of the beam energy in a target (see also

Sec. 12.2.2). The thickness of the anode region, in which the electron flux loses its energy, depends on time. Macroscopically, this fact is revealed in the strong nonstationarity of the anode impedance (especially of its reactive part). The dependence of both the characteristic thickness of a layer of the anode, in which the electron beam loses practically the whole energy, and, hence, the anode impedance on time is quite complex, since the layer thickness increases with a voltage on the diode (here, U is measured in V, and l_ε in cm):

$$l_\varepsilon(U) = \begin{cases} 2.1 \cdot 10^{-12} \frac{U^2}{\rho}, & 10 \text{ keV} < eU < 100 \text{ keV}, \\ 6.67 \cdot 10^{-11} \frac{U^{5/3}}{\rho}, & 100 \text{ keV} < eU < 1 \text{ MeV}, \\ \frac{5.11 \cdot 10^{-7} U - 0.26}{\rho}, & eU > 1 \text{ MeV}. \end{cases} \quad (12.188)$$

Moreover, l_ε decreases with increase in the current because of a growth of collective losses.

Owing to the development of nonlinear nonstationary processes, the transfers of energy and charge occur with a sharp front which moves with a finite velocity, and the mathematical model of the processes is reduced to evolutionary equations of the hyperbolic type, “telegraph” equations.

A radiophysical model of these processes involves delay lines leading to a delay of the pulses of current and voltage in the current guide relative to pulses incident on the anode surface, if the front of a current pulse incident on the surface is sufficiently short (as compared to the characteristic parameters of a line). Thus, in order to properly describe the electrophysical characteristics of a hard-current diode, we must add the anode represented by a delay line in its basic circuit (see Fig. 12.18). This element of the circuit realizes the limitation of the flux of electrons, the delay of a current pulse, and the accumulation of the excessive charge in the near-surface layer of the anode.

Collective losses of electrons in a dense plasma were studied in many works which yield that these effects are significant. The theoretical analysis of the correlative effects in the processes of losses (see, for example, Refs. 334–336) demonstrated a decrease of the characteristic length, on which the main share of the energy of electrons is lost, by several times. Upon the study of the energy losses by an electron flux in a substance, a state of the electron subsystem is assumed to be equilibrium. But the existence of a powerful flow of energy in the near-surface layer and of the flow regions in a substance with a strong disequilibrium (for example, a flow with the steep leading edge of a nonlinear wave) leads to the necessity to consider the fast processes of ionization of a target. These processes are the sources of a disequilibrium of the electron subsystem of a target (see Refs. 313–314).

In turn, the disequilibrium of electrons of the target further strengthens the processes of losses of the energy of electrons of the beam and provides a positive feedback (see Ref. 336). Nonequilibrium states of the electron subsystem significantly distort the profile of the losses of energy in a substance.

To properly estimate the losses of energy in a solid, it is necessary to determine the polarizability and conductivity under nonequilibrium conditions during the absorption of the energy of electrons. The inelastic interaction of electrons with a substance can be described in terms of dielectric permittivity $\tilde{\varepsilon}(q, \omega)$ (see Refs. 375, 376, 378) which is connected with the inelastic scattering cross-section by the relation

$$\frac{d^2\sigma}{d(\hbar\omega)d(\hbar q)} = \frac{1}{\pi a_0 q E} \text{Im} \left(-\frac{1}{\tilde{\varepsilon}(q, \omega)} \right). \quad (12.189)$$

Here E is the electron energy reckoned from the conduction band bottom, $\hbar\omega$ and $\hbar q$ are the changes of the energy and momentum of an electron as a result of the inelastic scattering, and $a_0 = \hbar^2/me^2$ is the Bohr radius. An analytic form of dielectric permittivity $\tilde{\varepsilon}(q, \omega)$ as a function of the frequency and wave vector is available only in some simple cases, for example, in an equilibrium state. In the collisions with the large transfer of energy and momentum, an electron behaves itself as a classical object. Therefore, $\omega_q \rightarrow \hbar q^2/2m$ as $q \rightarrow \infty$. The dispersion relations satisfying this condition can be chosen in various forms. For example, we can take the interpolational form Ref. 336,

$$\varepsilon^\ell(\omega, \vec{q}) = 1 + \frac{\omega_{ps}^2}{q^2 v_i^2} - \frac{\omega_0^2 + \frac{1}{2}\nu_F^2(\omega_0)q^2 + (\hbar q^2/2m)^2}{\omega_q^2} - i\pi f'(m\frac{\omega_q}{q}), \quad (12.190)$$

where ν_F the Fermi velocity. This approximation leads to the dispersion law which has the acoustic form (see Ref. 337) in the region of small wave numbers,

$$\omega_q^2 = \frac{\omega_0^2 + \frac{1}{2}\nu_F^2(\omega_0)q^2 + (\hbar q^2/2m)^2}{1 + \omega_{ps}^2/q^2 v_i^2}, \quad (12.191)$$

and passes to the classical limit in the region of large momenta transferred.

The distribution function of electrons over energy in a plasma-like medium is defined by a solution of the Boltzmann nonlinear equation with sources and sinks of particles (see Ref. 314) and has a power-like form. Simple estimates show that the energy losses upon the formation of such distributions increase by at least one order. For the analysis of a value of this nonequilibrium effect under real conditions of the experiment, it is necessary to solve a complicated inhomogeneous kinetic problem. However, upon the

injection of the electron beam in a target, there always exists a clear physical mechanism of limitation of the depth of penetration of the electron beam in a target. This mechanism gives the main contribution to the deceleration of the beam and can be analytically described, which will be demonstrated in what follows.

Such a mechanism is the accumulation of charge, creation of a virtual cathode, and stoppage of a beam upon the radial injection of a pulse hard-current electron beam in the cylindrical region occupied by a compensated plasma.

In the hydrodynamic approximation, the evolution of main physical quantities of the electron beam upon the radial injection of a pulse hard-current electron beam in the cylindrical region occupied by a dense plasma is described by the equations

$$\frac{\partial v}{\partial t} + v \frac{\partial v}{\partial r} = -\frac{e}{m} E - \frac{\gamma T}{m_e n_e} \frac{\partial n}{\partial z}, \quad (12.192)$$

$$\frac{\partial n}{\partial t} + \frac{1}{r^2} \frac{\partial}{\partial r} (r^2 n_e v_e) = 0, \quad (12.193)$$

$$\operatorname{div} \vec{E} = -4\pi e (n_e). \quad (12.194)$$

To solve these equations, we pass to the Lagrange variables $(r, t) \rightarrow (R, t_l)$, where R is the initial position of a Lagrange particle. The transition from the Euler variables to the Lagrange ones is realized with the help of the relations

$$\frac{\partial \psi}{\partial t_l} = \frac{\partial \psi}{\partial t} + \frac{\partial r}{\partial t} \frac{\partial \psi}{\partial r}, \quad \frac{\partial \psi}{\partial r} = \left(\frac{\partial \psi}{\partial R} \right) / \left(\frac{\partial r}{\partial R} \right), \quad \frac{\partial r}{\partial t_l} = u. \quad (12.195)$$

In the Lagrange variables, the equation of continuity Eq. 12.169 reads

$$\frac{\partial r}{\partial \tau} \cdot \frac{\partial (r^2 n_e)}{\partial t} + r^2 n_e \frac{\partial^2 r}{\partial t \partial \tau} = 0. \quad (12.196)$$

This equation can be integrated once, and its solution can be represented as

$$n_e = \frac{C_1(\tau)}{r^2 \partial r / \partial \tau}. \quad (12.197)$$

In the Lagrange variables with regard to Eq. 12.197, the Poisson equation Eq. 12.194 takes the form

$$\frac{\partial}{\partial \tau} (r^2 E) = -4\pi e r^2 \frac{\partial r}{\partial \tau} n_e(r) = -4\pi e \cdot C_1(\tau). \quad (12.198)$$

This equation can be also integrated, and its solution is

$$r^2 E = -4\pi e \int C_1(\tau) d\tau + \phi(t). \quad (12.199)$$

With regard to the derived solutions for the density and field, the equation of motion Eq. 12.192 written in the Lagrange variables becomes

$$\frac{\partial^2 r_2}{\partial t_l} = -\frac{e}{m} E - \frac{\gamma T}{m_e n_e} \frac{\partial n}{\partial \tau} \bigg/ \frac{\partial r}{\partial \tau} \frac{\partial r}{\partial \tau} - \frac{\gamma T}{m_c} \frac{r^2}{n_e \partial r / \partial \tau} \frac{\partial n}{\partial \tau}. \quad (12.200)$$

By multiplying the equation by $\frac{1}{r^2} (\partial r / \partial \tau)$ and integrating over τ , we arrive at the equation of motion

$$\begin{aligned} r \frac{\partial^2 r}{\partial t^2} - \int r \frac{\partial^3 r}{\partial t^2 \partial \tau} d\tau &= \frac{1}{r} \left[-\frac{4\pi e^2}{m} \left(\int C_1(\tau) d\tau - \phi(t) \right) \right] \\ &- \int \frac{1}{r} \left(-\frac{4\pi e^2}{m} \right) C_1(\tau) d\tau - \frac{\gamma T}{m} \ln n + \xi(t). \end{aligned} \quad (12.201)$$

The solution of this equation of motion is presented in Fig. 12.32. Near the maximum of the current in some small temporal region, the thickness of a layer, in which the energy is absorbed, decreases to a micron-scale size. This is connected to the fact that the penetration depth of electrons and, hence,

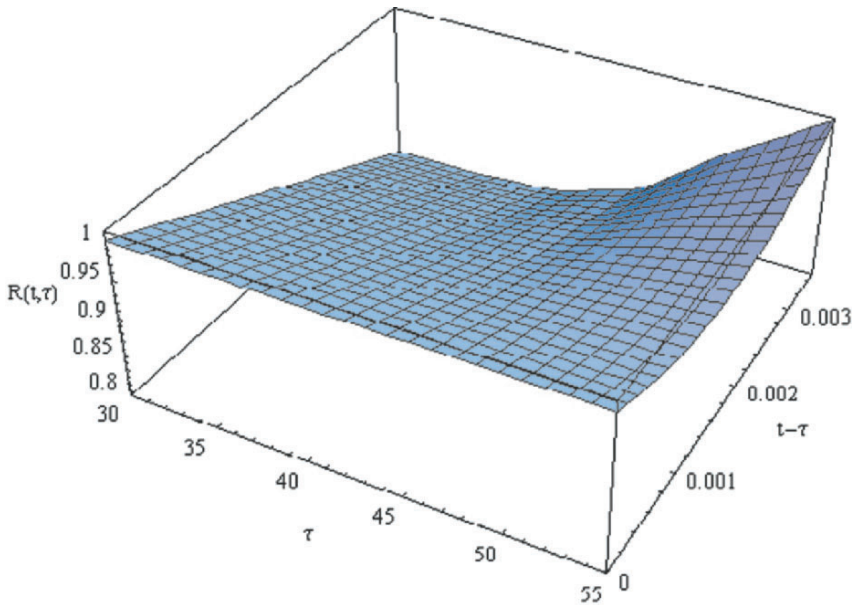


Fig. 12.32. Location of Lagrange particles for a flux of electrons as a function of the running time t and the injection time τ of electrons in the anode.

the region of losses of the energy of a REB firstly grow due to the increase in the energy of electrons and then rapidly decrease owing to the influence of collective effects on the loss of the energy of electrons in the anode.

We now construct a macroscopic model of the processes of transfer of electrons in a hard-current diode. The idea of a model consists in that the delay of a flux of electrons can be connected with increase in the reactive component of the impedance for the processes possessing a high-frequency component. In our process, a growth of the frequency is related to the steepening of a front.

For the short-pulse processes, the Maxwell equations can be reduced to a system of telegraph equations (see Ref. 338). Based on a qualitative description with simple hyperbolic equations, we now show how the frequency of a current pulse propagating in the system increases (the front becomes steeper).

We will seek a solution of the corresponding Cauchy problem for a simple wave as

$$v = F(t - r/u(v)). \quad (12.202)$$

In a solution of such a type, all dependent quantities can be expressed in terms of one of the quantities or through an auxiliary function $\Psi = \Psi(\tilde{r}, t)$. The last one satisfies the equation (see Ref. 339, 340)

$$\frac{\partial \psi}{\partial t} + \psi \frac{\partial \psi}{\partial \tilde{r}} = 0. \quad (12.203)$$

We assume that the relation $\psi = \psi_0$ is satisfied at $t = t_0$ and $\tilde{r} = r_0$. We represent ψ as $\psi = \psi_0 + \tilde{\psi}$ and introduce new variables $\xi = \tilde{\psi}/\psi_0$ and $r = \tilde{r}/\psi_0$.

Then the equation of evolution of perturbations of the charge density reads

$$\frac{\partial \xi}{\partial t} - (1 + \xi) \frac{\partial \xi}{\partial r} = 0. \quad (12.204)$$

This equation includes the parameters which are determined from the equation $\frac{dr}{dt} = 1 + \xi$. It is easy to see that

$$\frac{r - r_0}{1 + \xi} = t - t_0. \quad (12.205)$$

We assume that, at $r = r_0$, the initial distribution is defined by two characteristic times: the duration of a pulse and the duration of the leading edge. Then, upon the expansion of the initial distribution in a Fourier series

$$\Psi_0 = \sum_{i=1}^N A_i \sin(\omega_i t), \quad (12.206)$$

the basic frequencies will be connected just with these times, and the problem is reduced to solving the functional equation

$$\xi = \sum_{j=1}^N a_j \sin \left(\phi_j + \frac{q_j \xi}{1 + \xi} \right), \quad (12.207)$$

where a_j is the dimensionless amplitude of the j th signal at the input ($r = r_0$), q_j is the dimensionless coordinate, and ϕ_j is the dimensionless phase. This equation defines the spectrum of the combination frequencies of oscillations of the density which appear in the system upon the propagation of a perturbation along the coordinate r .

The functional equation was solved in Ref. 340 under a quite weak restriction imposed on amplitudes ($(a_1 + a_2) < 1$). It is convenient to represent the sine in Eq. 12.207 through the exponential by the Euler formula and to write the functional equation as

$$\xi = \sum_{i=1}^N \frac{a_j}{2i} \left[\exp \left(i \left(\phi_j + \frac{q_j \xi}{1 + \xi} \right) \right) - \exp \left(-i \left(\phi_j + \frac{q_j \xi}{1 + \xi} \right) \right) \right]. \quad (12.208)$$

The further transformations of the equation are based on the representation of the exponential as a series in generalized Laguerre polynomials $(1 - z)^{-\alpha-1} \exp \left(\frac{xz}{z-1} \right) = \sum_{n=0}^{\infty} L_n^{\alpha}(x) z^n$, $|z| < 1$. As a result of awkward transformations of the series, we get the exact solution of the functional equation Eq. 12.208. The asymptotic series for this solution reads

$$\xi = 2 \sum_{l=0}^{\infty} \sum_{p=0}^{\infty} \sum_{\substack{j=1 \\ l+p \neq 0}}^N (-1)^{jp} \frac{1}{Q_{lp}} J_l(a_1 Q_{ep}) J_p(a_2 Q_{lp}) \sin \phi_{lp}, \quad (12.209)$$

where $Q_{lp} \equiv Q = lq_1 + (-1)^j p q_2$ is the combination argument defining the amplitudes of waves, $\phi_{lp} = l\phi_1 + (-1)^j p \phi_2$ is the combination phase. In the case where only one wave is excited, i.e., at $p = 0$ and $a_2 = 0$, we get the classical formula of Bessel–Fubini

$$\xi = 2a_1 \sum_{l=1}^{\infty} \frac{J_e(la_1 q_1)}{la_1 q_1} \sin l\phi_1. \quad (12.210)$$

Thus, the above-presented expression generalizes formula Eq. 12.210 to the case of interacting waves (see Refs. 340, 342).

Eq. 12.209 is symmetric relative to the permutation of indices, which allows us to consider Q as a positive value in the further calculations.

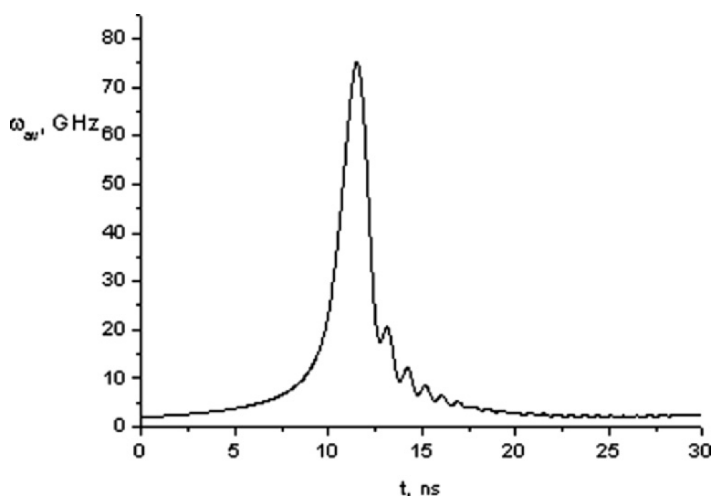


Fig. 12.33. Evolution of effective frequency of a current pulse as the result of the sharpening of its front.

At $l = 0$, we should not sum over j . The case where $p = 0$ is derived from the previous one by a simple change of indices.

Thus, we derive explicitly the solution of the problem of two interacting waves. In the case where $l = 1$ and $p = 0$ ($l = 0, p = 1$), we get the amplitudes of basic frequencies as functions of z . Such a dependence for various values of z is presented in Fig. 12.33.

The case where $l \neq 0$ and $p \neq 0$ yields the spectrum of combination frequencies, where $j = 1$ corresponds to “red” satellites, and $j = 2$ to “violet” ones.

The derived solution has a completed form, where the expression under the symbol of summation represents the appropriate partial contribution of harmonics to the spectrum.

The analysis of the solution yields that it is valid up to the overturning of a nonlinear wave, which occurs when $a_1 q_1 + a_2 q_2$ is equal approximately to 1. It is seen that the spectrum becomes significantly wider while approaching the time moment of a steepening of the wave: the spectrum includes very high harmonics with considerable amplitudes.

According to such an evolution of the frequency of perturbations of the density of the flux of electrons, the diode impedance grows sharply with characteristic frequencies of the current. For the typical parameters of an anode and voltage pulses, the dependence of impedance on time calculated by this model is shown in Fig. 12.34, and the dependence of the thickness λ_ε of a layer, in which the beam energy is absorbed, on time is given in Fig. 12.35.

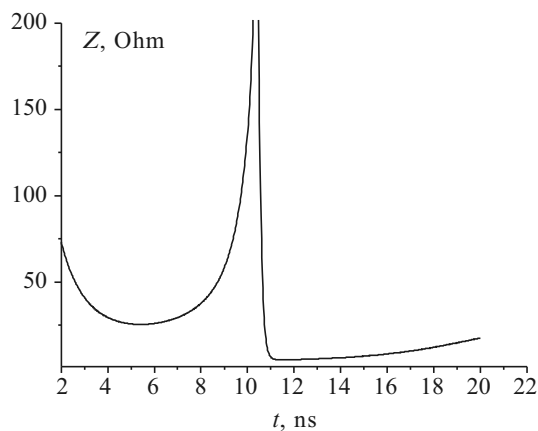


Fig. 12.34. Typical dependence of the anode impedance on time.

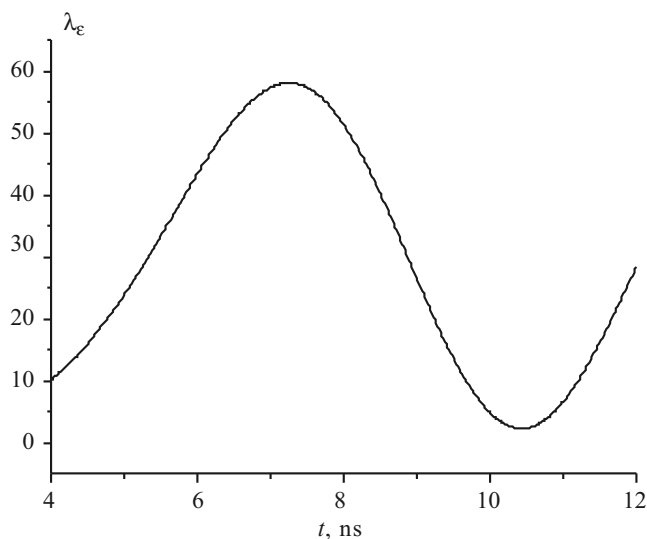


Fig. 12.35. Dependence of the near-surface layer thickness, in which the main share of the beam energy is absorbed, on time.

As a consequence of such peculiarities of the electrodynamic properties of the diode, the density of absorbed energy in a near-surface layer of the anode sharply grows. In Fig. 12.36, we give the estimate of the pressure pulse form on the anode surface which arises due to the collective processes providing the high density of the energy released in the transient layer between the dense plasma flare and a target.

The figures demonstrate that the collective processes in the anodic plasma and a near-surface layer of the anode induce an acute self-consistent

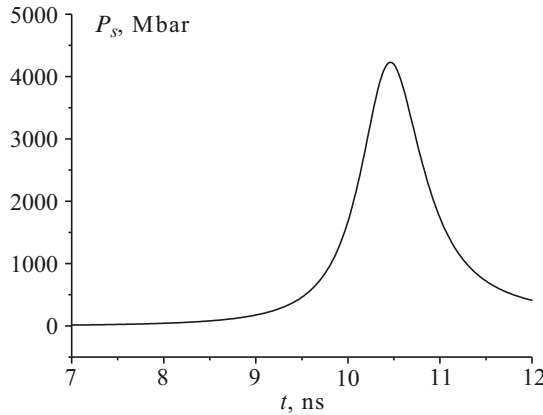


Fig. 12.36. Typical dependence of the pressure in a near-surface layer of the anode on time.

sharpening of a pressure pulse on the surface anode. The energy flow density in the near-surface layer of a target and the pressure in this region increase with time by a law corresponding to the mode with sharpening.

In the following sections, we consider how the arising mode with sharpening excites a nonlinear wave, whereas the electron charge being accumulated with time exceeds the fluctuation level of the density (see Refs. 341, 342) and forms a dynamical virtual cathode which creates, during its evolution, a nonlinear wave of volume charge.

We emphasize that the described situation is possible not always, but only upon the use of a thoroughly designed diode with regard to peculiarities of the dynamics of a beam of electrons, a plasma cathode, and a plasma anode.

Hydrodynamic Equations of a Target with the Adiabatic Equation of State under a Pulse Action on the Surface. In this section, we consider the excitation of radial hydrodynamic flows (nonlinear waves) arising as a result of the development of the above-described processes with sharpening in a thin near-surface layer of the anode and propagating along a radius.

In cylindrical coordinates, hydrodynamic equations are usually written in the form of a system of equations including the equation of continuity

$$\frac{\partial \rho}{\partial t} + \frac{1}{r} \frac{\partial}{\partial r} (r \rho u) = 0; \quad (12.211)$$

the equation of motion of a substance in the radial direction

$$\frac{\partial u}{\partial t} + u \frac{\partial u}{\partial r} = -\frac{1}{\rho} \frac{\partial P}{\partial r} + a_s(t) D_a (r - R); \quad (12.212)$$

and the evolutionary equation for the entropy

$$\frac{\partial S(P, \rho)}{\partial t} + u \frac{\partial S(P, \rho)}{\partial r} = 0. \quad (12.213)$$

Instead of Eq. 12.213, we can use equations presented in various forms which are convenient in various physical situations. In particular, if the adiabatic equation of state holds, Eq. 12.213 can be written in a simple form as

$$\frac{\partial}{\partial t} \left(\frac{P}{\rho^\gamma} \right) + u \frac{\partial}{\partial r} \left(\frac{P}{\rho^\gamma} \right) = 0. \quad (12.214)$$

For ultra-fast processes, it is convenient sometimes to replace Eq. 12.207 by the energy conservation equation and the evolutionary equation for the energy flow. We present these equations in the form of the equations for temperature and a heat flow. The evolution of temperature is defined by the equations following from the energy conservation law with regard to a source of energy concentrated in the near-surface layer,

$$C_0 \rho \frac{\partial T}{\partial t} + \frac{w}{r} + \frac{\partial w}{\partial z} = Q_s(t) D_a(r - R), \quad (12.215)$$

and the equation for a flow of energy

$$\tau(T) \frac{\partial w}{\partial t} + w + k(T) \frac{\partial T}{\partial r} = 0. \quad (12.216)$$

By introducing a new variable $n_\rho(r, t) = r\rho(r, t)$ and Lagrange variables $(r, t) \rightarrow (R, t_l)$ (see Eq. 12.195), the equation of continuity Eq. 12.211 can be easily integrated as above with regard to Eq. 12.197.

Consider a version with the adiabatic equation of state. In this case, the equation for entropy in the Lagrange variables can be written as $\frac{\partial}{\partial t_l} \left(\frac{P}{\rho^{\gamma_a}} \right) = 0$, and its solution has the simple form: $P/\rho^{\gamma_a} = C_s(R) \equiv P_0(R)/\rho_0(R)^{\gamma_a}$.

Substituting this solution in the equation of motion, we get

$$\begin{aligned} \frac{\partial^2 r}{\partial t^2} = & -\frac{r}{\rho_0(R)R} \left(\frac{\partial C_s(R)}{\partial R} \rho^{\gamma_a} + C_s(R) \gamma_a \rho^{\gamma_a-1} \frac{\partial \rho}{\partial R} \right) \\ & + a_s(t) D_a(r - R), \end{aligned} \quad (12.217)$$

where the right-hand side includes an external force, whose typical form is $a_s(t) = P_0/[t_{ex}^\alpha (1 - t/t_{ex})^\alpha]$ in the self-consistent mode.

Let us arrange Lagrange particles at the initial time moment along the system axis and consider the evolution of this set of Lagrange particles

R_i . For each particle, the equation of motion has the form Eq. 12.217 with the initial values R_i . We denote the density of a target at points R_i at the initial time moment as ρ_{oi} . Then the density of a target at the next time moments is defined by the solution of the equation of continuity in the form Eq. 12.197:

$$\rho_i(R_i, t) = \frac{\rho_0(R_i) R_i}{r_i(R_i, t) \partial r_i(R_i, t) / \partial R_i}. \quad (12.218)$$

To solve the equations of motion, it is necessary to set a specific model of external action. We take into account that the external action is concentrated only on the surface and rapidly decreases inwards the substance. In the analytic studies, we may take $D_a(r - R_0) = \delta(r - R_0)$, and, in the numerical analysis, we will use one of the continuous approximations of the delta-function:

$$D_a(x) = \frac{1}{\pi} \frac{\varepsilon}{x^2 + \varepsilon^2}, \quad \varepsilon \ll 1. \quad (12.219)$$

Let us also write the equations in the case where we use the energy conservation equation (the heat conduction equation). For the equation of state, we take a simple approximation with the use of the Grüneisen coefficient γ ($\gamma_a \approx \gamma$) as $P = \gamma_a \rho T$ which is frequently employed in the analysis of the action of a pulse hard-current beam of electrons on an anode (see Ref. 236). For copper, $\gamma_a = 0.875\gamma$ and $\gamma = 1.91$.

We now write the complete system of equations in Lagrange variables with the use of the general relations Eq. 12.195. The dynamical equations for N_i Lagrange particles of the medium are

$$\frac{\partial r_i}{\partial t} = u_i(R_i, t), \quad (12.220)$$

$$\frac{\partial u_i(R_i, t)}{\partial t} = -\beta \frac{1}{\rho_i(R_i, t)} \left(\left(\frac{\partial P_i(R_i, t)}{\partial R_i} \right) / \left(\frac{\partial r_i(R_i, t)}{\partial R_i} \right) \right), \quad (12.221)$$

$$\rho_i(R_i, t) = \frac{\rho_0(R_i) R_i}{r_i(R_i, t) \partial r_i(R_i, t) / \partial R_i}, \quad (12.222)$$

$$\begin{aligned} \tau \left(\frac{\partial W_i(R_i, t)}{\partial t} - \frac{\partial r_i(R_i, t)}{\partial t} \left(\frac{\partial W_i(R_i, t)}{\partial R_i} \right) / \left(\frac{\partial r_i(R_i, t)}{\partial R_i} \right) \right) \\ + W_i(R_i, t) + \kappa \left(\frac{\partial T_i(R_i, t)}{\partial R_i} \right) / \left(\frac{\partial r_i(R_i, t)}{\partial R_i} \right) = 0, \end{aligned} \quad (12.223)$$

$$\rho_i(R_i, t)C_v \left(\frac{\partial T_i(R_i, t)}{\partial t} - \frac{\partial r_i(R_i, t)}{\partial t} \left(\frac{\partial T(R_i, t)_i}{\partial R_i} \right) / \left(\frac{\partial r(R_i, t)}{\partial R_i} \right) \right) + \left(\frac{\partial W_i(R_i, t)}{\partial R_i} \right) / \left(\frac{\partial r_i(R_i, t)}{\partial R_i} \right) + \frac{W(R_i, t)_i}{r_i} = Q(r_i(R_i, t), t); \quad (12.224)$$

$$i = 1, \dots, N_l. \quad (12.225)$$

In the chosen units, the coefficient $\beta = 4.28$.

To solve the system of equations of the dynamics of Lagrange particles, it is convenient to introduce the deviation $\xi(R_i, t)$ of the i -th Lagrange particle from the initial position:

$$r_i(R_i, t) = R_i - \xi(R_i, t). \quad (12.226)$$

For deviations, we get the system of nonlinear wave equations with a disturbing force $F(\xi_i R_i, t)$:

$$\frac{\partial^2 \xi_i(R_i, t)}{\partial t^2} - C_d^2(\xi_i R_i, t) \frac{\partial^2 \xi_i(R_i, t)}{\partial R_i^2} = F(\xi_i R_i, t), \quad (12.227)$$

where the squared velocity of propagation of a perturbation is defined by the relation

$$C_d^2(\xi_i R_i, t) = \frac{\gamma_s T_i(R_i, t)}{(1 - \partial \xi_i / \partial R_i)^2}, \quad (12.228)$$

and the force $F(\xi_i R_i, t)$ acting on a Lagrange particle consists of the force from the side of the medium $F_{\text{int}}(\xi_i, R_i, t)$ and the external force (pressure) $a_s(t)D_a(r_i(R_i, t) - R_0)$ developing in the external layer of a target:

$$F(\xi_i R_i, t) = F_{\text{int}}(\xi_i, R_i, t) - a_s(t)D_a(r_i(R_i, t) - R_0). \quad (12.229)$$

The pressure $F_{\text{int}}(\xi_i, R_i, t)$ depends on a medium state and is defined by the formula

$$F_{\text{int}}(\xi_i, R_i, t) = -\gamma_s T_i(R_i, t) \times \left(\frac{1}{R_i - \xi_i(R_i, t)} - \frac{1}{R_i (1 - \partial \xi_i(R_i, t) / \partial R_i)} - \frac{1}{\rho_0(R_i)} \frac{\partial \rho_0(R_i)}{\partial R_i} \right) + \frac{\gamma_s}{1 - \partial \xi_i(R_i, t) / \partial R_i} \frac{\partial T_i(R_i, t)}{\partial R_i}. \quad (12.230)$$

The system of equations Eq. 12.227 describes a nonlinear wave which is excited on the surface and moves to the center. The propagation velocity of

the wave depends on the external force and can reach great values exceeding the sound velocity in the equilibrium state.

The analysis of the data of Kiev experiments on the transmutation of elements shows (under the assumption that the nuclear transformations in the anode occur in the region where the extremum parameters are reached in a target) that the steepening of a wave, i.e., an increase in the wave amplitude and a decrease in the characteristic thickness of its front, can run at distances significantly less than the radius of a target R_0 .

We will show how such states appear as solutions of the system of hydrodynamical equations in Lagrange variables with regard to the fact that the formation of a nonlinear wave occurs on a periphery. For this purpose, we introduce new Lagrange variables b_i instead of the Lagrange variables R_i by the relation $R_i/R_0 = 1 - b_i$. If the evolutionary processes run mainly on a periphery, then the quantities b_i are small ($b_i \ll 1$). In this case, the actual coordinates of a particle $r_i(R_i, t)$ are connected with b_i by the relation

$$r_i(R_i, t)/R_0 = 1 - b_i - \xi_i(b_i, t). \quad (12.231)$$

In first approximation, we will consider the distribution of temperatures to be prescribed.

To numerically solve the system of equations Eq. 12.227, it is necessary to know the functions $\partial \xi_i(b_i, t)/\partial b_i$ and $\partial^2 \xi_i(b_i, t)/\partial b_i^2$ at every time moment. By using the smallness of $b_i \ll 1$, it is convenient to represent a solution of the system of equations as a series

$$\xi_i(b_i, t) = \sum_{k=0}^{\infty} f_k(t) b_i^k \quad (12.232)$$

and to get a simple approximation for derivatives.

Substituting these expansions in the system of equations with partial derivatives Eq. 12.227 and equating the coefficients of the identical powers of b_i , we get the unclosed system of ordinary differential equations for the expansion coefficients $f_k(t)$:

$$f_0'' = \gamma_s T \left(-\frac{1}{1 - f_0(t)} - \frac{1}{1 + f_1(t)} + \frac{2f_2(t)}{(1 + f_1(t))^2} \right) - \frac{\varepsilon a_s(t) K}{\pi (\varepsilon^2 + f_0(t)^2)}, \quad (12.233)$$

$$f_1'' = \gamma_s T \left(-\frac{1+f_1(t)}{(1-f_0(t))^2} - \frac{1}{1+f_1(t)} + \frac{2f_2(t)}{(1+f_1(t))^2} - \frac{8f_2^2(t)}{(1+f_1(t))^3} \right) - \frac{\varepsilon (2f_0(t)(1+f_1(t))) a_s(t) K}{\pi (\varepsilon^2 + f_0(t)^2)^2}, \quad (12.234)$$

$$f_2'' = \gamma_s T \left(-\frac{1+f_1(t)}{(1-f_0(t))^2} - \frac{1}{1+f_1(t)} + \frac{2f_2(t)}{(1+f_1(t))^2} - \frac{8f_2^2(t)}{(1+f_1(t))^3} \right) + \gamma_s T \left(-\frac{1+f_1(t)}{(1-f_0(t))^2} - \frac{1}{1+f_1(t)} + \frac{2f_2(t)}{(1+f_1(t))^2} - \frac{8f_2^2(t)}{(1+f_1(t))^3} \right) - \frac{\varepsilon (2f_0(t)(1+f_1(t))) a_s(t) K}{\pi (\varepsilon^2 + f_0(t)^2)^2}. \quad (12.235)$$

It is essential that, after some transformations of the first integrals of the starting equations Eqs. 12.221–12.225, we can prove that all the coefficients $f_k(t)$ are interconnected between themselves by recurrence algebraic relations expressing $f_k(t)$ in terms of the previous coefficients and the external force. Thus, these recurrence relations reduce the problem to a single strongly nonlinear and inhomogeneous equation for $f_0(t)$. The first recurrence relations are as follows:

$$f_1(t) = -1 - \frac{\exp(-a_s(t)K)}{1-f_0(t)}, \quad (12.236)$$

$$\frac{f_2(t)}{(1+f_1(t))^2} = \frac{1}{4} \left(a_s(t)K + \frac{f_1(t) + 2f_0(t)}{(1-f_0(t))(1+f_1(t))} \right). \quad (12.237)$$

To numerically solve the system of equations Eqs. 12.221–12.225, we will use the approximation of the external force which follows from its estimates based on the above-discussed model of a diode (see Fig. 12.18). The model function used in the numerical analysis is given in Fig. 12.37.

By truncating the chain of equations and restricting themselves by a finite number of terms of the expansion, we can find all the hydrodynamic variable we are interesting in.

The results of the numerical analysis of the system of equations are given in Figs. 12.38–12.42.

The figures show the spatio-temporal evolution of the substance density for various intensities of the external source. At small intensities, we observe the appearance of a high-density region in the near-surface layer of a target as a result of the action of the source (Fig. 12.38).

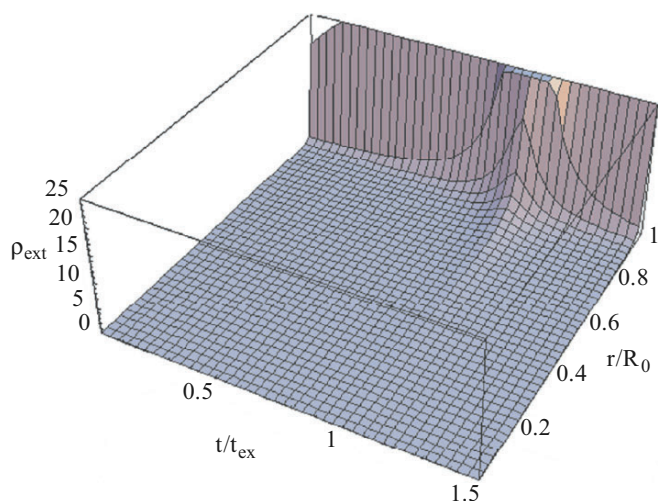


Fig. 12.37. External force versus Lagrange variables (in order to show the spatial structure, values of the functions are cut at a relatively small level).

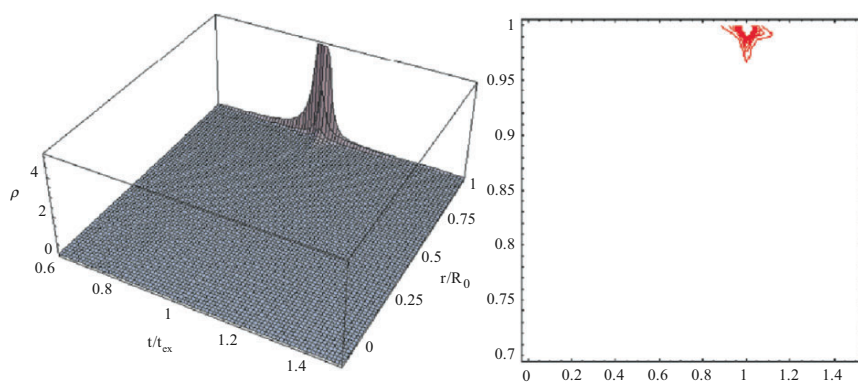


Fig. 12.38. Evolution of the substance density at small intensities of the action (the dimensionless amplitude of the action is taken to be 0.01).

With increase in the source intensity, the features of the evolution of the density become pronounced more clearly, as seen in Figs. 12.39–12.40. A very narrow high-density region is originated in the near-surface layer. Located between close radii, this region represents itself as a thin wave shell (a wave-shell) moving to the center. The shell is a nonlinear wave of the pulse form rushing in the substance. In connection with that the basic equations of motion have the form of the equations for nonlinear oscillators, we can observe a spatio-temporal structure during the movement of the shell.

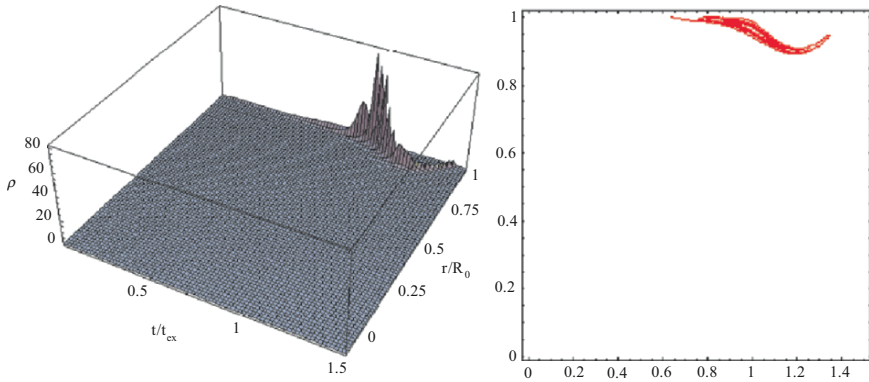


Fig. 12.39. Evolution of the substance density at moderate intensities of the action (the dimensionless amplitude of the action is taken to be 0.05).

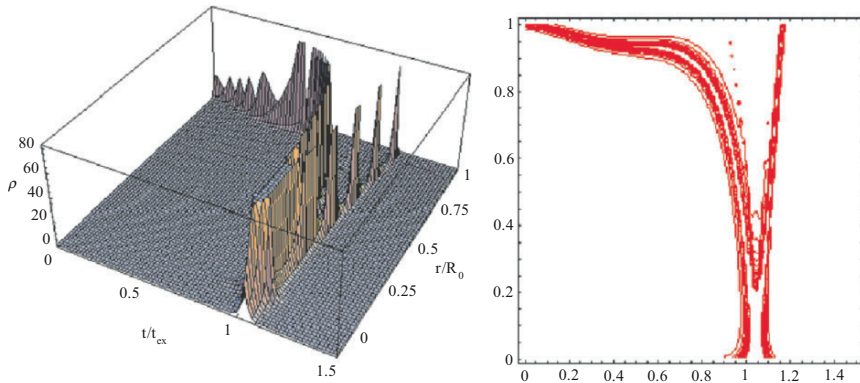


Fig. 12.40. Evolution of the substance density at a great intensity of the action (the dimensionless amplitude is of the order of 1.5). Density amplitudes are cut for the clearness. Maximum values of the peak amplitudes are equal to several thousands.

In Fig. 12.41, we show the evolution for a great intensity of the action in the form of the spatial profile of a wave at various time moments.

The derived solutions can be approximated analytically for the one-fluid ground state as

$$u_{e0}(r, t) = \frac{u_{\max}}{\sqrt{1 + \sigma^2(r - r_0(t))^2}}, \quad (12.238)$$

$$n_{e0}(r, t) = C_e \sigma \sqrt{1 + 1/\sigma^2(r - r_0(t))^2}. \quad (12.239)$$

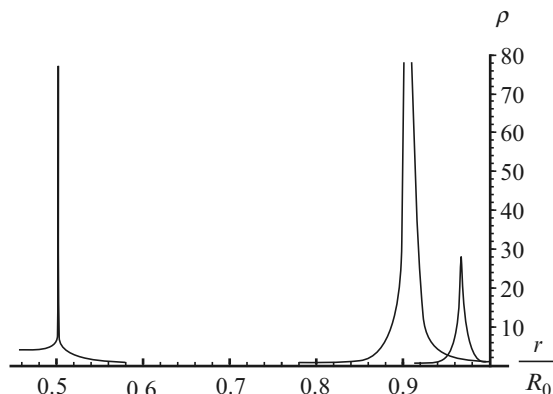


Fig. 12.41. Distributions of the substance density in a target along a radius at various time moments ($t/t_{ex} = 0.2$, $t/t_{ex} = 0.7$, and $t/t_{ex} = 0.95$).

Here, $r_0(t)$ is the trajectory of the region of sharpening of the non-linear wave.

At a large energy flow from an external source, a wave arises in a very narrow spatial region near the surface and propagates along a radius (see Fig. 12.39).

Under a pulse action with sufficiently large intensity on a continuum, shock waves are usually excited in it. Formally, shock waves arise due to the appearance of the ambiguity upon the crossing of the trajectories of Lagrange particles. The excitation of shock waves hampers the compression of a substance to high densities. As shown in Introduction, the mode with sharpening allows one, at least in principle, to get any density of a substance and to avoid the appearance of shock waves (discontinuities).

In Fig. 12.42, we show the trajectories of close Lagrange particles near the pulse maximum in the case where the mode with sharpening is absent.

It is seen that the trajectories cross one another. Upon the crossing, a discontinuity or a shock wave is formed.

Figure 12.43 demonstrates the trajectories of Lagrange particles in the case of the mode with sharpening.

It is seen that the trajectories approach one another very closely, but do not intersect. Just this permanent approach of particles ensures the permanent increase in the substance density during a realization of the mode with sharpening.

The dynamics of Lagrange particles of a substance has some peculiarities. In Figs. 12.44 and 12.45, we display the velocity and acceleration of two particles from the near-surface layer. It is seen that, up to the time moment

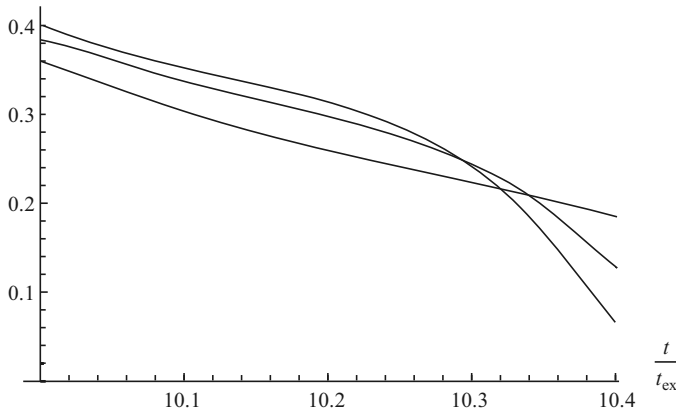


Fig. 12.42. Trajectories of three close Lagrange particles near the pulse maximum without the mode with sharpening.

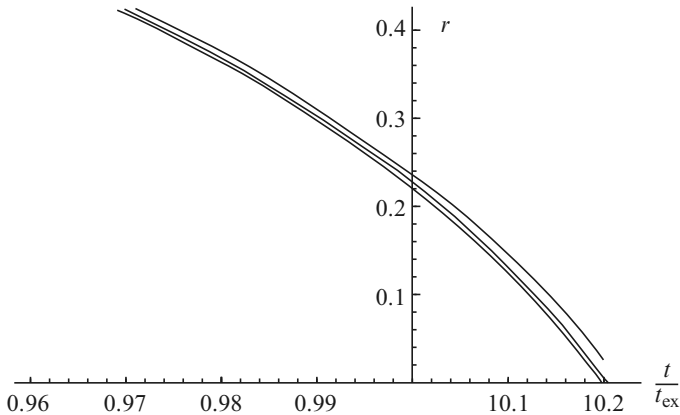


Fig. 12.43. Trajectories of three close Lagrange particles near the time moment when the external pressure sharpening occurs in the case of a high-intensity flow.

of the order of $0.55t_{ex}$, a steady movement from the near-surface layer to the center is formed as a result of the competition of the external force and the resistance of the medium. Due to the fast increase in the external action intensity, there exists a time moment ($0.55t_{ex}$), after which the external action constantly exceeds the resistance force. After this time moment in the mode with sharpening, we observe the appearance of both the acute growth of the acceleration of the high-density region to the center and the increase in the density. Such a dynamics leading to an acute increase in gradients is a basis for the appearance of instabilities of the gradient type in a dense plasma, the separation of charges, and the appearance of self-consistent electric fields on the front of a nonlinear wave moving to the center.

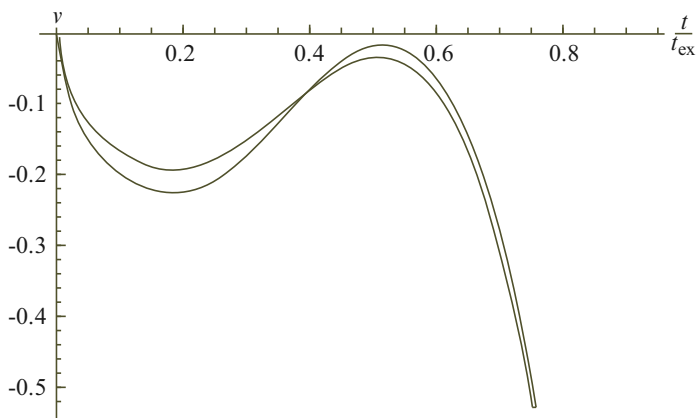


Fig. 12.44. Velocity of Lagrange particles versus time.

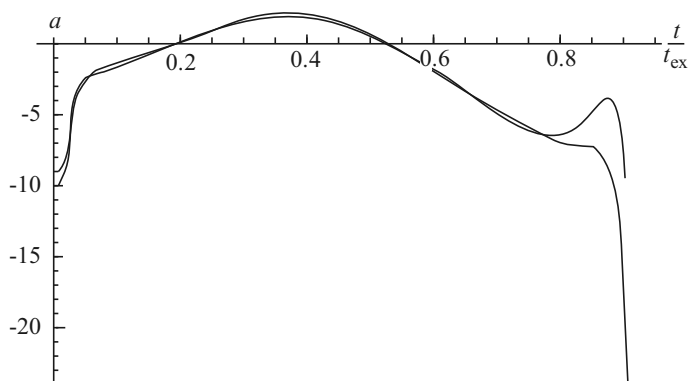


Fig. 12.45. Acceleration of Lagrange particles versus time.

In the next section, we will consider such instabilities which induce the appearance of the two-fluid hydrodynamic mode of evolution of a substance under the action of a pulse electron beam.

12.2.4. Instabilities of the One-Fluid Flow and the Excitation of a Two-Fluid Flow of the Electron-Nucleus Plasma

The evolution of a one-fluid flow of the target substance can pass in the two-fluid stage under the action of a pressure pulse on the surface due to the development of an instability in the one-fluid flow when some parameters reach critical values.

Figure 12.39 shows the possible sequence of the processes running in the anode, being under the action of a pulse REB. Depending on the

character of the nonstationarity of REB and its maximum current, the squeezing of a target can be realized in several modes:

- in the one-fluid mode;
- in the one-fluid mode with sharpening;
- in the one-fluid mode with sharpening and with the development of the instability ensuring the transition into the two-fluid mode and its nonlinear stage, namely a plasma-field structure;
- in the mode with sharpening passing in the two-fluid stage due to a rapid growth of the volume charge delivered by the beam into the target with the subsequent transition into a strongly nonlinear stage, namely the formation of a plasma-field structure.

The direct transition from the mode with sharpening in the two-fluid stage is considered in the next section. Here, we continue the analysis of a one-fluid flow and study the excitation instability of the two-fluid mode. As a result of the development of this instability, the densities of electrons and ions stop to satisfy the condition of quasineutrality, and there arises the two-fluid stage of the evolution. Because the maximum values of the target parameters are attained in the region of a superdense shell arisen in the one-fluid mode, just the hydrodynamic state near the shell should be considered as the zero approximation upon the analysis of the instability.

For the analysis of the processes in a target, it is extraordinarily important to know the value of the increment of this instability and the dependence of the increment on the main parameters of the target, pressure pulse, and shell. Due to the great difficulties of the analysis of instabilities of a strongly inhomogeneous plasma, we below consider a simple variant allowing us to analyze the increment under assumption of the preset evolution of the target temperature.

We now write the system of equations of the two-fluid hydrodynamics of a one-component plasma for a 1D cylindrically symmetric flow of ion and electron fluids with regard to self-consistent electric fields as follows: the equations for the ion flow,

$$\frac{\partial u_{ir}}{\partial t} + u_{ir} \frac{\partial u_{ir}}{\partial r} - \frac{u_{i\theta^2}}{r} = -\frac{Ze}{m_i} E_r - \omega_{Bi} u_{i\theta} - R_{ir} - \frac{1}{m_i n_i} \frac{\partial P_i}{\partial r}, \quad (12.240)$$

$$\frac{\partial u_{i\theta}}{\partial t} + u_{i\theta} \frac{\partial u_{i\theta}}{\partial r} + \frac{u_{i\theta} u_{ir}}{r} = \omega_{Bi} u_{ir} - R_{i\theta}; \quad (12.241)$$

the continuity equation for the ion flow,

$$\frac{\partial n_i}{\partial t} + \frac{1}{r} \frac{\partial}{\partial r} (r n_i u_{ir}) = 0; \quad (12.242)$$

the equations for the electron flow,

$$\frac{\partial u_{er}}{\partial t} + u_{er} \frac{\partial u_{er}}{\partial r} - \frac{u_{e\theta}^2}{r} = -\frac{e}{m} E_r - \omega_{Be} u_{e\theta} + R_{er} - \frac{1}{m_e n_e} \frac{\partial P_e}{\partial r}, \quad (12.243)$$

$$\frac{\partial u_{e\theta}}{\partial t} + u_{e\theta} \frac{\partial u_{e\theta}}{\partial r} + \frac{u_{e\theta} u_{er}}{r} = \omega_{Be} u_{er} + R_{e\theta}; \quad (12.244)$$

the continuity equation for the electron flow,

$$\frac{\partial n_e}{\partial t} + \frac{1}{r} \frac{\partial}{\partial r} (r n_e u_{er}) = 0, \quad (12.245)$$

and the equation for a self-consistent electric field,

$$\frac{1}{r} \frac{\partial}{\partial r} (r E_r) = 4\pi e (n_e - Z n_i). \quad (12.246)$$

Here, $R_{er} = \nu_{ei} (u_i - u_e)$ and $R_{ir} = \nu_{ei} (u_i - u_e)$ are the friction forces acting on the electron and ion flows, respectively.

The basic one-fluid state characterized by $u_{e0}(r, t)$, $n_{e0}(r, t)$, $u_{i0}(r, t) = u_{e0}(r, t)$, and $n_{i0}(r, t) = n_{e0}(r, t)/Z$ has been studied in the previous section. For the analysis of the stability of this state, it is necessary to analyze the evolution of a deviation from it:

$$u_e(r, t) = u_{e0}(r, t) + \tilde{u}_e(r, t), \quad n_e(r, t) = n_{e0}(r, t) + \tilde{n}_e(r, t), \quad (12.247)$$

$$u_i(r, t) = u_{i0}(r, t) + \tilde{u}_i(r, t), \quad n_i(r, t) = n_{i0}(r, t) + \tilde{n}_i(r, t), \quad (12.248)$$

etc.

From the initial system Eqs. 12.240–12.246, we get the following equations for the perturbations of densities and velocities of the electron and ion flows in the radial direction taking the first approximation:

$$\begin{aligned} m_e n_{e0} \left[\frac{\partial \tilde{v}_e}{\partial t} + \frac{\partial (v_{e0} \tilde{v}_e)}{\partial r} + \frac{e}{m_e} \tilde{E} \right] + m_e \tilde{n}_e \left[v_{e0} \frac{\partial v_{e0}}{\partial r} \right] \\ = -\lambda_{ei} (\tilde{v}_e - \tilde{v}_i) - \gamma_e T_e \frac{\partial \tilde{n}_e}{\partial r}, \end{aligned} \quad (12.249)$$

$$\begin{aligned} M_i n_{i0} \left[\frac{\partial \tilde{v}_i}{\partial t} + \frac{\partial (v_{i0} \tilde{v}_i)}{\partial r} + \frac{Ze}{M} \tilde{E} \right] + M_i \tilde{n}_i \left[v_{i0} \frac{\partial v_{i0}}{\partial r} \right] \\ = -\lambda_{ei} (\tilde{v}_i - \tilde{v}_e) - \gamma_i T_i \frac{\partial \tilde{n}_i}{\partial r}, \end{aligned} \quad (12.250)$$

$$\frac{\partial \tilde{n}_e}{\partial t} + \frac{1}{r} \frac{\partial}{\partial r} [r (n_{e0} \tilde{v}_e + \tilde{n}_e v_{e0})] = 0, \quad (12.251)$$

$$\frac{\partial \tilde{n}_i}{\partial t} + \frac{1}{r} \frac{\partial}{\partial r} [r (n_{i0} \tilde{v}_i + \tilde{n}_i v_{i0})] = 0, \quad (12.252)$$

$$\frac{1}{r} \frac{\partial}{\partial r} (r \tilde{E}) = 4\pi e (Z \tilde{n}_i - \tilde{n}_e). \quad (12.253)$$

On the linear stage of the development of the instability, we assume that $\tilde{n}_i \approx 0$ and derive the simple relations for the perturbations of the density and velocity of the electron flow:

$$\tilde{n}_e = -\frac{1}{4\pi e} \left(\frac{1}{r} \frac{\partial}{\partial r} (r E_r) \right), \quad (12.254)$$

$$\tilde{v}_e = \frac{1}{n_{e0}} \left(\frac{1}{4\pi e} \frac{\partial}{\partial t} (E_r) + \frac{C(t)}{r} - v_{e0} \tilde{n}_e \right). \quad (12.255)$$

Substituting these relations in the equation for the velocity of a perturbation of the electron flow, we get the equation for the self-consistent radial electric field in plasma:

$$\begin{aligned} & \frac{\partial^2 E_r(r, t)}{\partial t^2} - 2v_{e0}(r, t) \frac{\partial^2 E_r(r, t)}{\partial r \partial t} + V^2(r, t) \frac{\partial^2 E_r(r, t)}{\partial r^2} + \Omega^2(r, t) E_r(r, t) \\ & = a(r, t) \frac{\partial E_r(r, t)}{\partial r} + b(r, t) \frac{\partial E_r(r, t)}{\partial t} + \psi(r, t). \end{aligned} \quad (12.256)$$

All coefficients in this wave equation for a perturbation of the electric field have obvious physical sense and are expressed through the velocity and density of the electron flow in the one-fluid approximation. The perturbation propagation velocity is defined by both the velocity of the unperturbed flow and the temperature:

$$V^2(r, t) = v_{e0}^2(r, t) - \gamma T(r, t), \quad (12.257)$$

$$\begin{aligned} \Omega^2(r, t) = \omega_{pe}^2 & + \left[\frac{\gamma T}{r^2} + \frac{\nu_{eff} v_{e0}}{r} - \frac{v_{e0}^2}{r^2} - \frac{v_{e0}}{r} \frac{\partial_t n_0}{n_0} + \frac{\partial_t v_{e0}}{r} \right. \\ & \left. + \frac{(\gamma T - v_{e0}^2)}{r} \frac{\partial_r n_0}{n_0} + \frac{2v_{e0} \partial_t v_{e0}}{r} \right], \end{aligned} \quad (12.258)$$

$$\begin{aligned} a(r, t) = \nu_{eff} v_{e0} & - \frac{\gamma T}{r} + \frac{v_{e0}^2}{r} - v_{e0} \frac{\partial_t n_0}{n_0} + \partial_t v_{e0} \\ & + (\gamma T - v_{e0}^2) \frac{\partial_r n_0}{n_0} + 2v_{e0} \partial_r v_{e0}, \end{aligned} \quad (12.259)$$

$$b(r, t) = \nu_{eff} + \frac{v_{e0}}{r} - \frac{\partial_t n_0}{n_0} - v_{e0} \frac{\partial_r n_0}{n_0} + \partial_r v_{e0}. \quad (12.260)$$

The exciting force $\psi(r, t)$ depends on the parameters of the main flow and on the form of the total current passing through the system:

$$\begin{aligned} \psi(r, t) = & \frac{4\pi e \partial_t C(t)}{r} + \frac{4\pi e C(t)}{r} \\ & \times \left(\nu_{eff} - \frac{v_{e0}}{r} - \frac{\partial_t n_0}{n_0} - v_{e0} \frac{\partial_r n_0}{n_0} + \partial_r v_{e0} \right). \end{aligned} \quad (12.261)$$

The total current enters into the expression for the force via the integration constant $C(t)$.

As the zero approximation, we will use the approximations of the solutions for a one-fluid flow which have been derived in the previous section. In Figs. 12.46–12.49, we give the main characteristics of a nonlinear density wave having the form of a moving shell. A trajectory of the density maximum of a shell upon its movement to the center is shown in Fig. 12.46.

Upon the movement, the shell becomes thinner, and its density increases. The density maximum and the shell thickness versus time are given, respectively, in Figs. 12.47 and 12.48.

Upon the propagation of a nonlinear wave over the anode, all the macroscopic characteristics of the substance are changed. In the analysis of the instability, we must consider both the density and the distributions of the velocity and temperature, whose characteristic forms are shown, respectively, in Figs. 12.49 and 12.50.

Since the expression $V^2(r, t) - v_{e0}^2(r, t) = \gamma T(r, t)$ in Eq. 12.256 for the self-consistent field is positive [$\gamma T(r, t) > 0$], the equation is of the hyperbolic type in the whole region of variations of the variables. By virtue of the hyperbolicity, Eq. 12.256 has real characteristics which are defined by the equation

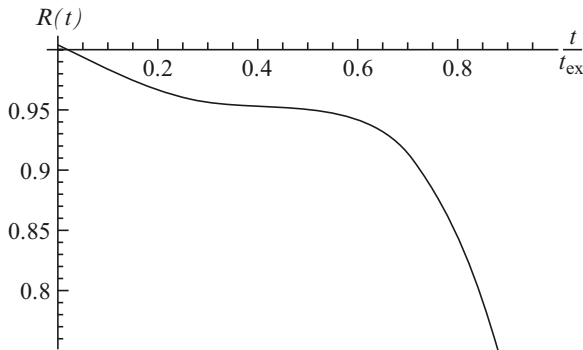


Fig. 12.46. Trajectory of the density maximum.

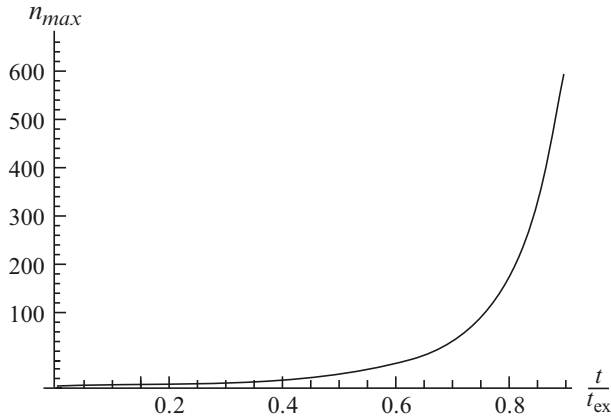


Fig. 12.47. Density maximum versus time.

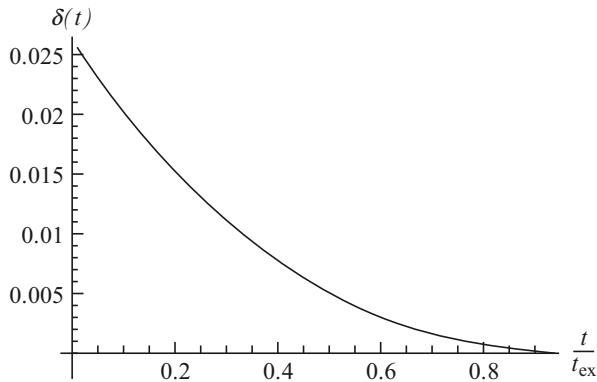


Fig. 12.48. Shell thickness versus time.

$$\frac{dr(t, \tau)}{dt} = v(r, t) - \gamma T(r, t), \quad r(0, \tau) = 1, \quad (12.262)$$

where τ is the time moment of the start of the evolution of a characteristic from the anode surface, and the fields of velocities and temperature have a characteristic form shown in Figs. 12.49 and 12.50.

Let us estimate the instability of the appearance of a self-consistent field on the characteristics of the equation for a nonlinear wave, by representing the field of the wave in the form of a Gauss pulse moving along the characteristic $r = R(t, \tau)$ with width $\delta(t)$ and amplitude $A(t)$:

$$E_r(r, t) = A(t) \exp \left[- \left(\frac{r - R(t, \tau)}{\delta(t)} \right)^2 \right]. \quad (12.263)$$

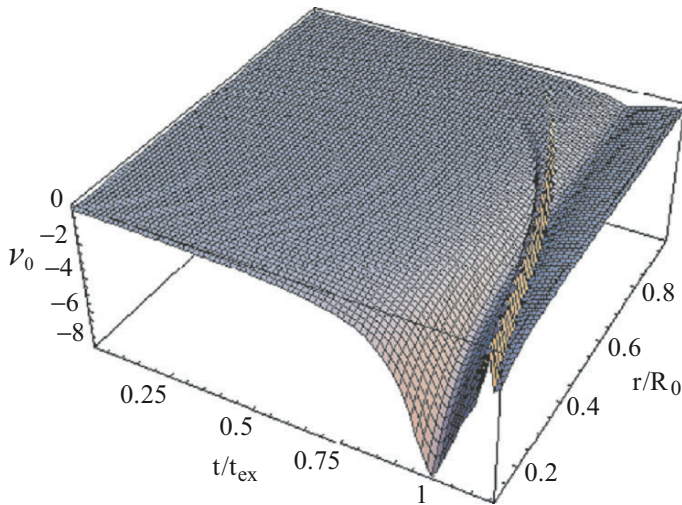


Fig. 12.49. Evolution of the velocity distribution for a one-fluid flow along the radius.

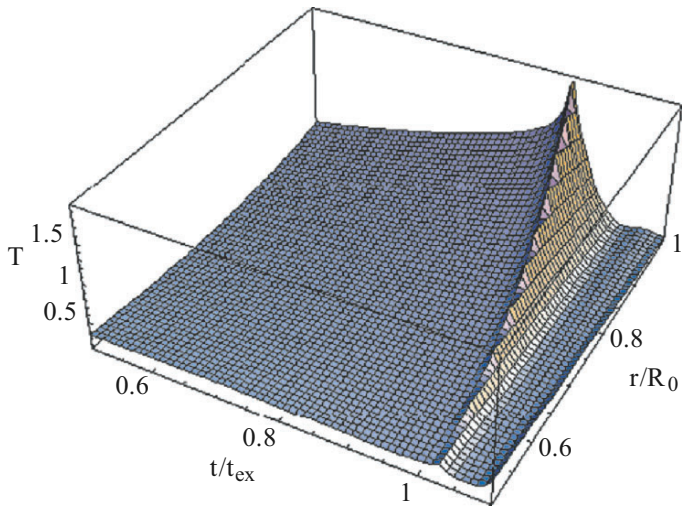


Fig. 12.50. Evolution of the temperature distribution for a one-fluid flow along the radius.

Equation 12.256 for the electric field intensity yields the equation for the amplitude

$$\frac{d^2 A}{dt^2} + 2g(t)\frac{dA}{dt} + \Omega^2(t)A(t) = F(t), \quad (12.264)$$

which is similar to the equation for an oscillator excited by an external force $F(t)$.

The typical forms of the temporal dependencies of the frequency, damping coefficient $g(t)$, and exciting force $F(t)$, which were calculated for the one-fluid mode, are presented in Figs. 12.51–12.53.

The nonstationarity of the coefficients in the equation for an oscillator leads to the development of instabilities with large increment. Moreover, the wave amplitude for the self-consistent field increases from the level of fluctuations to very great values in agreement with the field equation Eqs. 12.262–12.264.

The evolution of distribution of electron density corresponding to the development of this instability is given in Fig. 12.54.

The growth of the self-consistent field in the dense plasma means, in fact, the transition of the system into the two-fluid mode. In this case, there

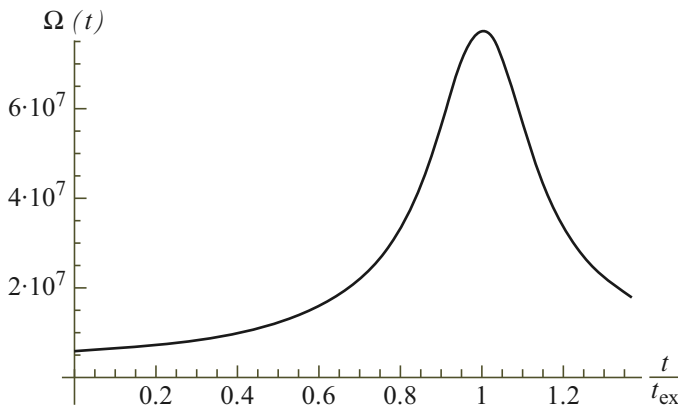


Fig. 12.51. $\Omega(t)$ versus time.

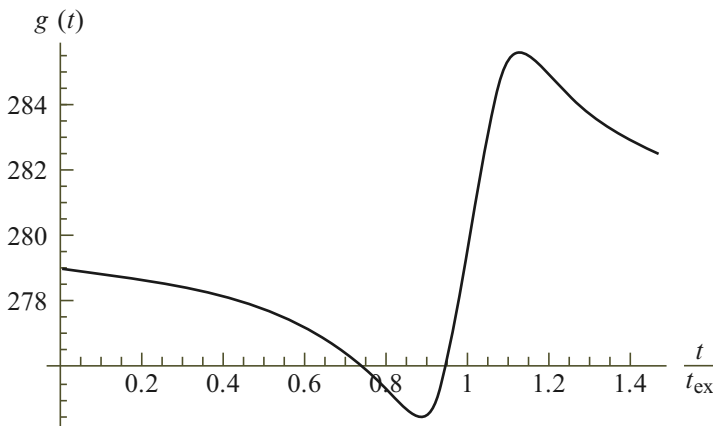


Fig. 12.52. $g(t)$ versus time.

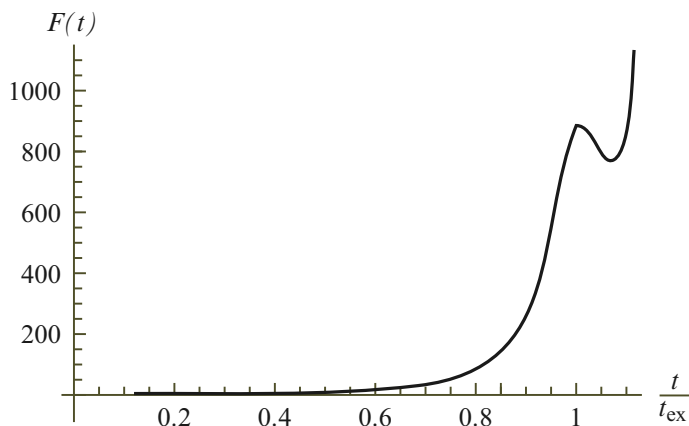


Fig. 12.53. Exciting force $F(t)$ versus time.

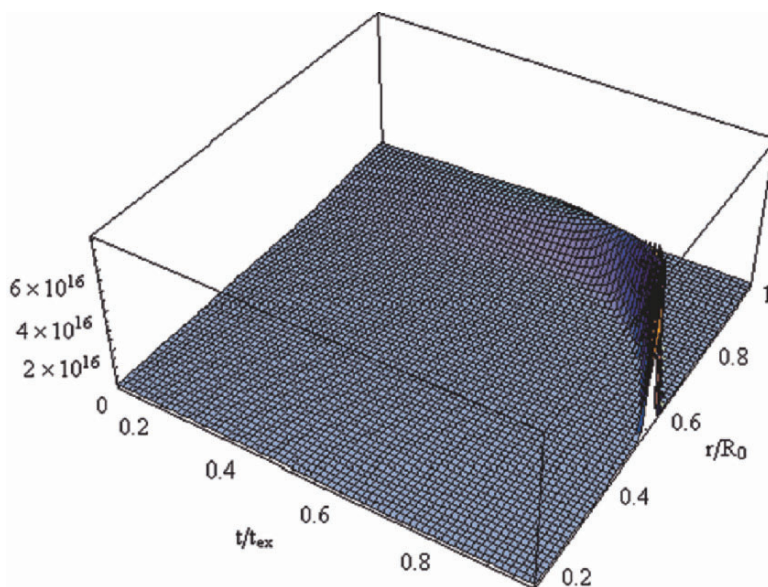


Fig. 12.54. Typical evolution of the self-consistent distribution of electron densities.

occurs the separation of charges, and the self-consistent field is propagating in the form of a nonlinear wave of the volume charge to the anode center. On the well-developed nonlinear stage, this wave is a plasma-field structure which will be described in the subsequent sections.

12.2.5. Two-Fluid Mode upon the Action of a Pulse Electron Beam on a Target

The two-fluid mode of a plasma flow can arise, as mentioned at the beginning of this section (see also Fig. 12.39), not only as a result of the transition of a one-fluid flow in this mode due to the development of an instability, but also at the very beginning of the process of interaction of the beam of electrons with the target. In the last case, the two-fluid mode arises as a result of the accumulation of the excessive charge, whose value exceeds the level of fluctuations in the dense plasma and can become very great in the course of the nonlinear evolution.

The main collective effects in plasma are related mostly to the excitation of self-consistent fields. Self-consistent electric fields in plasma reach great amplitudes usually as a result of the regular linear and nonlinear plasma oscillations. In Ref. 342, we have demonstrated the new possibility to derive extremely high fields in a dense nonideal plasma due to the development of nonlinear processes which are similar to those leading to high densities of plasma in the experimental studies of both the generation of neutrons in a deuterium mixture in electrostatic fields (Ref. 343) and the electrostatic or magneto-electric confinement of plasma (Ref. 344).

As will be shown below, the self-consistent description of electrodynamic and hydrodynamic processes allows one to define the conditions for the appearance of plasma-field structures in the dense plasma of a target, the structures being composed from virtual cathodes and anodes. The characteristic spatial scales of a structure decrease from the periphery to the center or to the system axis. Such a reduction of scales leads to a significant increase in the field and its inhomogeneity.

The effects related to the appearance of the structures of virtual electrodes are so far essential for the whole series of applications that, already from the middle of the last century, the special national programs aimed at the study of the physics of the relevant processes and at their practical applications are realized.

As shown at the beginning of Sec. 12.2.3, the specific conditions of the action of a pulse REB on the anode in the experiments performed at the Electrodynamics Laboratory "Proton-21" induce the limitation of the flow of electrons in the surface layer on the macroscopic level, which is manifested in the specific evolution of the diode impedance shown in Fig. 12.34. Such a temporal behavior of the impedance leads to a delay of both the current pulse in the cathode-anode gap $I_c(t)$ and the current in an inverse conductor $I_b(t)$. This delay is well seen in Figs. 12.55 and 12.56.

The difference of the currents $[I_c(t) - I_b(t)]$ induces the accumulation of electrons in a thin layer, in which the beam is decelerated and loses its

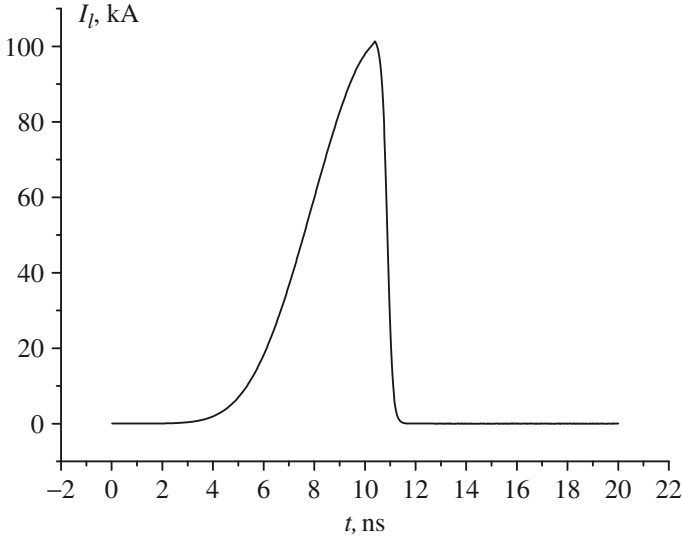


Fig. 12.55. Current in the cathode-anode gap.

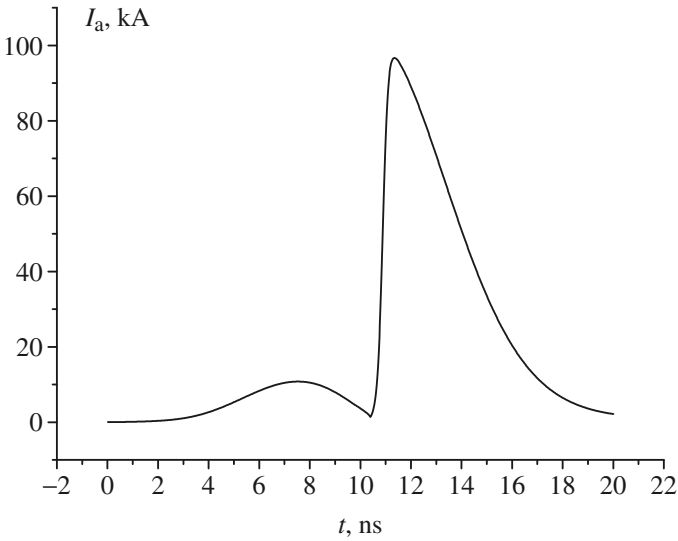


Fig. 12.56. Current in the inverse conductor.

energy. In Fig. 12.57, we present the evolution of the numbers of ions and electrons in this layer.

Such an accumulation of the volume charge leads to the necessity to consider further the hydrodynamic processes with regard to self-consistent electric fields in the two-fluid approximation.

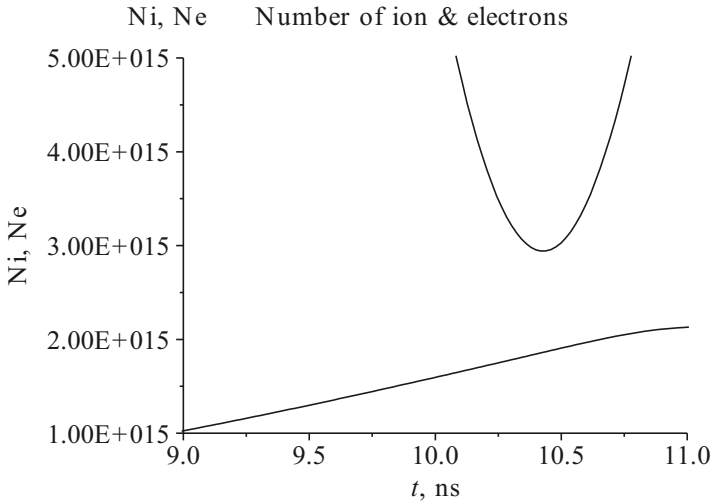


Fig. 12.57. Numbers of electrons delivered by the beam (lower curve) and ions (electrons) of the metal in the surface layer (upper curve) versus time.

Interaction of a Pulse Hard-Current Electron Beam with a Target in the Two-Fluid Mode. In the framework of the two-fluid hydrodynamics of a one-component plasma with regard to self-consistent fields in the cylindrically symmetric case, the system of Eqs. 12.240–12.246 for the flows of the ion and electron fluids includes the equations of the ion flow, the equation of continuity for the ion flow, the equations of the electron flow, the equation of continuity for the electron flow, and the equation for a self-consistent electric field.

Upon the action of electron beams on the condensed medium, the electron subsystem of a target reacts most rapidly. Ion flows propagate with lower velocities than those of electron ones and can be considered as transport flows for the electron flows arising against their background.

In what follows, we will analyze the equations for conduction electrons in a metal in the case of the cylindrical symmetry of a target undergoing the external action. We set the distribution of ions $n_i = n_{i0}(r)$ as that defined by the one-fluid mode with sharpening.

Like in the case of the one-fluid mode, we will solve the hydrodynamic equations by using the Lagrange variables.

In the previous case, the system consisted of a substance filling all the region $(0, R_0)$ which was subjected to the external action, and we studied the dynamics of particles of this substance in the one-fluid approximation. That is, the substance was considered to be quasineutral, and its dynamics was considered as that of neutral masses. In that case, the Lagrange variables

were, naturally, the initial coordinates of particles of the medium and time. Now the situation is more complex, because our system consists of electrons which are injected on the target surface at $r = R_0$ at various time moments τ (see Sec. 12.2.3 and of ions which are distributed over the whole volume of the target together with the target electrons neutralizing them. Therefore, it is necessary to use the initial coordinates and time as the Lagrange variables for ions (see Sec. 12.2.3) and the time moments of the injection and time for the beam electrons. Let us now rewrite the total derivative $d/dt \equiv \partial/\partial t + u_{er}\partial/\partial r$, which is given in the Euler variables, in the Lagrange ones (τ, t_l) , where τ is the time moment of the entry of Lagrange particles, being the beam electrons, into the solid-state plasma. Then $u_{er} = \partial r/\partial t_l$, and the solution of the equation of continuity can be presented, as above, in the form

$$rn_e \frac{dr}{d\tau} = C_1(\tau). \quad (12.265)$$

It is easy to see that $C_1(\tau)$ is defined by both the injection current $I_c(t)$ of electrons from the cathode and the target surface area $S_a(t)$, on which the electron beam falls: $C_1(t) = \alpha_c I_c(t)/S_a(t)$.

Let us pass to the Poisson equation. By multiplying both sides of it by $r \frac{\partial r}{\partial \tau}$, we substitute the solution of Eq. 12.274, integrate one time, and get the expression for the field as

$$rE_r = 4\pi e \left[\int_0^\tau \alpha_c I_c(\tau_1) d\tau_1 - \int_0^\tau n_{i0}(\tau_1) r \frac{\partial r}{\partial \tau_1} d\tau_1 \right] + \chi(t). \quad (12.266)$$

The function $\chi(t)$ in Eq. 12.265 can be determined from the analysis of the Maxwell equations with curl and is related to the total current passing through the system, i.e., to the current in the inverse conductor $I_b(t)$. Thus, we have $\chi(t) = \alpha_b I_b(t)$.

The equation for the angular velocity can be also integrated; and, as a result, we obtain

$$u_{e\theta} = \frac{\omega_{Be}}{2} r + \frac{C_4(\tau)}{r}. \quad (12.267)$$

Upon the injection of electrons in the target, both the axial component of the current and the transverse one $I_\perp(t) = \alpha_\perp I_c(t)$ arise at the expense of collective processes. Just the last component of the electron injection current defines the constant $C_4(t) = \frac{\omega_{Be}(I_\perp(t))}{2} R_0^2$ (for the sake of simplicity, we consider the case of the zero initial azimuth velocity of the beam upon its fall onto the target).

By using the presented analytic solutions, we arrive at the equation of motion for a Lagrange particle,

$$\begin{aligned} r \frac{d^2 r}{dt^2} + \nu r \frac{dr}{dt} - \frac{\omega_{Be}(I_{\perp}(t))}{2} \frac{R_0^2}{r^2} + \frac{\omega_{Be}^2}{4} r^2 \\ = -\frac{4\pi e^2}{m_e} \int_0^{\tau} \alpha_c \frac{I_c(\tau_1)}{S_a(\tau_1)} d\tau_1 + \frac{4\pi e^2 n_i}{m_e} \frac{r^2}{2} - \frac{4\pi e^2}{m_e} \alpha_b I_b(t). \end{aligned} \quad (12.268)$$

Consider the evolution of a set of Lagrange particles with the time moments τ_i of their entry into the system. The coordinate of the entry into the system is R_0 for all Lagrange particles. By n_{0ei} , we denote the density of electrons in the target at the initial time moments τ_i . The density of electrons in the target at the subsequent time moments is defined by the formula

$$n_{ei}(\tau_i, t) = \frac{n_{0ei}(\tau_i) R_0}{r_i(\tau_i, t) \partial r_i(\tau_i, t) / \partial \tau_i}. \quad (12.269)$$

By using this solution of the equation of continuity, we can write the system of equations describing the dynamics of Lagrange particles as

$$\frac{\partial r_i}{\partial t} = u_{ei}(\tau_i, t), \quad (12.270)$$

$$\begin{aligned} r_i \frac{du_{ei}}{dt} + \nu r_i u_{ei} - \frac{\omega_{Be}^2(I_{\perp}(t)) R_0^4}{4r_i^2} + \frac{\omega_{Be}^2}{4} r_i^2 \\ = -\frac{4\pi e^2}{m_e} \int_0^{\tau_i} \alpha_c \frac{I_c(\tau_1)}{S_a(\tau_1)} d\tau_1 + \frac{4\pi e^2 n_i(r_i)}{m_e} \frac{r_i^2}{2} - \frac{4\pi e^2}{m_e} \alpha_b I_b(t). \end{aligned} \quad (12.271)$$

This system of ordinary differential equations for Lagrange particles was solved numerically with the use of some models of the processes running in the diode which are described at the beginning of this chapter.

The phase portrait of dynamic system of Lagrange particles representing the result of the numerical integration of the system of equations is given in Fig. 12.58.

The dynamics of Lagrange particles presented in this figure induces the excitation of a nonlinear pulse of the excessive charge and, hence, the self-consistent field. The density wave of the volume charge, which propagates to the center of the system, becomes steeper, and the self-consistent field increases strongly. The last process is accompanied by the flows of electrons and ions crossing the region of its steepening. The evolution of the wave of the volume charge passes then into its final stage, namely in the formation of plasma-field structures.

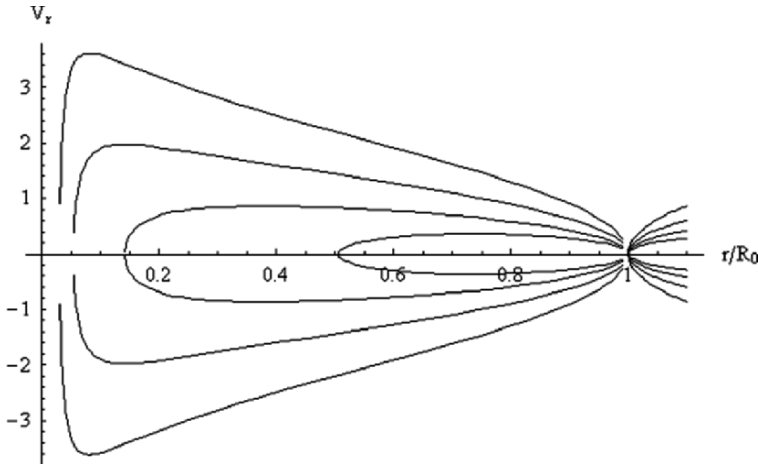


Fig. 12.58. The phase portrait of a dynamic system of Lagrange particles representing the nonlinear plasma wave steepening.

12.2.6. Structures in the Electron-Nucleus Plasma and a Mechanism of the Energy Transportation onto the Nuclear Scale

As above, we consider the 1D models of the evolution of the system along the radius. The 1D models can possess the cylindrical or spherical symmetry. In this case, the cylindrically symmetric models seem to be more natural under the conditions realized in our experiments. However, the cylindrically symmetric squeezing of the target passes, in some cases, into the spherically symmetric one.

We explain this in the following way. The cylindrically symmetric passage of the current along the anode leads to states which are similar to the pinch states of beams in plasma. During the evolution, these states become unstable and induce the perturbations which are periodic along the current propagation direction (see Fig. 12.59).

In the presence of the additional symmetrizing factor related to the cumulative effect (see Fig. 12.59), there arises an inhomogeneity in the growth of the amplitude of the unstable oscillation. This inhomogeneity causes the primary development of one of the semiperiods of the oscillation. The described evolution terminates in making the process to be more spherically symmetric.

In the case of the developing processes with sharpening, the formation of a cylindrically symmetric core of the target possessing a sufficiently high temperature become possible. Here, we may expect the appearance of an instability of the Rayleigh–Taylor type. The evolution of this instability will

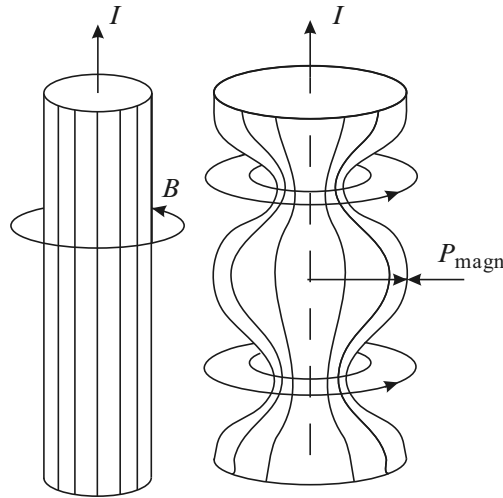


Fig. 12.59. Appearance of the instability upon the propagation of the current.

be realized by a close scenario. The estimates of the characteristic scales of the instability for typical parameters of the beam give the symmetrization region size $R_c \sim 50 \mu\text{m}$.

A steepening of the spatial-charge wave, which has been described in the previous section, leads to a sharp growth of the electric field and to the appearance of the flows of electrons and ions. The further evolution of nonlinear waves is realized via the formation of plasma-field structures.

For the structure of fields in plasma-beam systems, essential are the peculiarities of the joint dynamics of the electron and ion (nuclear) flows. For example, the structure of fields is significantly affected by the hydrodynamic flow of ions, which interacts with the electron relativistic flow possessing a chaotic (thermal) component (see Ref. 345). In many cases, for example in plasma-beam discharges, it is necessary to consider the thermal component for both electrons and ions. We note that the type of an originating plasma-field structure depends on the dominance of the hydrodynamic and/or kinetic contribution to the dynamics of charged particles.

In particular, the evolution of macroscopic parameters (temperature, current, pressure, etc.) admits the appearance of the systems of double layers of the ion-acoustic, electron-acoustic, or magneto-acoustic nature (see Ref. 341).

In some cases, these double layers can be a system of virtual electrodes formed against the background of the modes with sharpening, as indicated at the beginning of this section. The attainment of a sufficient density and the separation of charges (the appearance of a self-consistent electric field)

in thin double layers can create the conditions for the union of nuclei in clusters due to collisions.

In Ref. 347 and related ones, the processes of formation of structures in plasma were analyzed with the use of the kinetic approximation for ions and electrons. The formation of structures in the hydrodynamic approximation was considered in Refs. 345, 346.

We now present the consequences of the dynamics of structures in the radial direction (see Ref. 346). We will consider a spherical region of the anode occupied by the electron-nucleus plasma, where the flows of electrons and nuclei are moving.

Let us analyze both the flow of positively charged nuclei with the thermal component and the electron beam, for which we neglect temperature-related corrections. Quasistationary states of the electron-nucleus plasma are described by the system of equations including the equations of continuity,

$$\operatorname{div}(n_e \vec{v}_e) = 0, \quad \operatorname{div}(n_i \vec{v}_i) = 0, \quad (12.272)$$

and the equations of motion for electrons and nuclei in the hydrodynamic approximation. If we consider only radial motions, these equations of motion can be integrated once, and we get the energy preservation equations

$$\frac{m_e}{2} \vec{v}_e^2 - e\varphi = \frac{m_e}{2} v_0^2; \quad \frac{m_i}{2} \vec{v}_i^2 + Ze\varphi = T_i. \quad (12.273)$$

The system is closed by the field equation

$$\Delta\varphi = 4\pi e (n_e(\varphi) - Z n_i(\varphi)). \quad (12.274)$$

In the structures with spherical symmetry, Eqs. 12.272–12.274 possess peculiarities leading to a sharper reduction of the scale of the potential near the center as compared to the case of cylindrical symmetry.

In this case, the nonlinear equation for the potential reads

$$\frac{d}{dr} \left(r^2 \frac{d\Psi}{dr} \right) = \frac{\beta_e}{\sqrt{1+\Psi}} - \frac{\beta_i}{\sqrt{1-\eta\Psi}}, \quad (12.275)$$

where $\Psi = 2e\varphi/m_e v_{e0}^2$, $\eta = Zm_e v_{e0}^2/v_{i0}^2$, $\beta_e = 2eI_e/m_e v_{e0}^2 l_a v_{e0}$, $\beta_i = 2eI_i/m_i v_{i0}^2 l_a v_{i0}$, I_e , I_i are the electron and ion currents, respectively, and l_a is the anode length.

The solution of the equation is presented in Fig. 12.60.

For the central region where $r \approx 0$ and $\Psi \approx -1 + y$, this equation is reduced to the Emden–Fowler equation

$$\frac{d}{dr} \left(r^2 \frac{dy}{dr} \right) = \beta_e y(r)^{-1/2}, \quad (12.276)$$

which can be analyzed in detail (see Ref. 348).

At $r \rightarrow 0$, the asymptotics has the form $y(r) = ar^b$. According to the general theorems (see Ref. 348), the characteristic scales of oscillations decrease while approaching the center. It is seen from Fig. 12.60 that the structure of virtual electrodes has the properties of similarity (being a fractal structure).

At the expense of the reduction of scales and the increase of the amplitude of oscillations while approaching the system axis, the electric fields in plasma can become very intense, as seen in Fig. 12.61.

The densities of the electron and ion components also form structures with scales reducing toward the center. The structures of the electron and ion components and that of the potential are somewhat shifted each relative

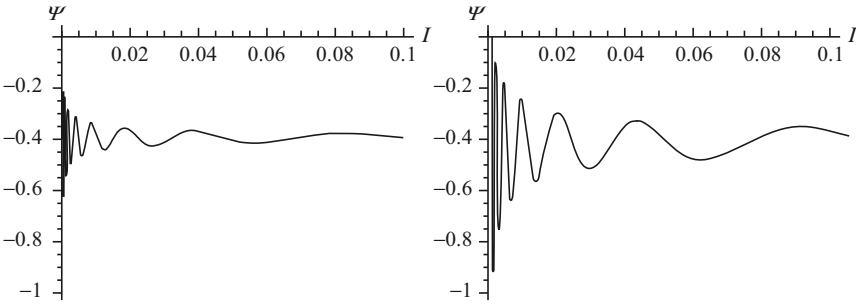


Fig. 12.60. Structure of the potential in relative units in a spherical region of plasma with virtual electrodes.

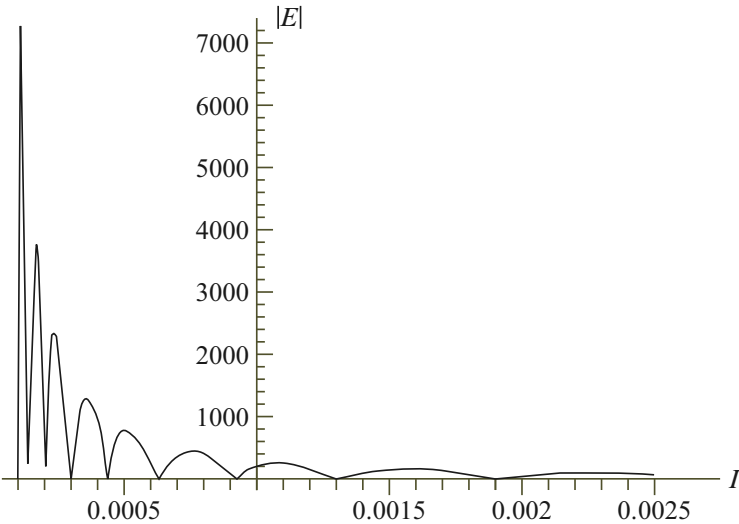


Fig. 12.61. Dependence of the electric field on the radius in a spherical region of plasma with virtual electrodes.

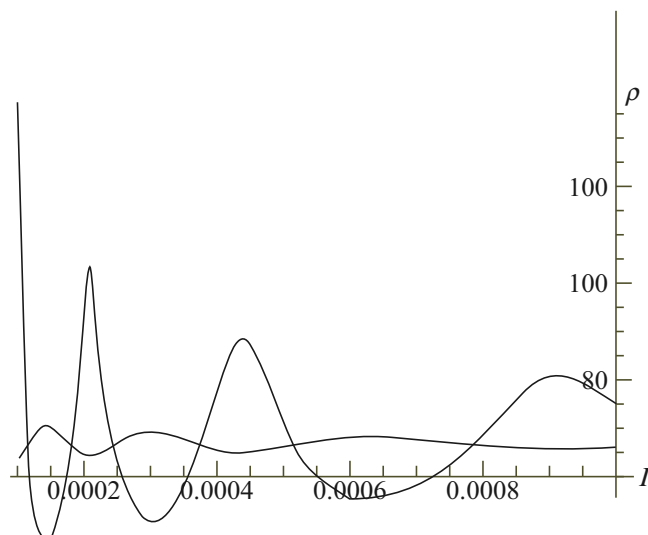


Fig. 12.62. Structure of the density of ions and electrons in relative units in a spherical region of plasma with virtual electrodes.

to other by phase. As seen from Figs. 12.61 and 12.62, the ion component density and the field amplitude increase, respectively, by at least 2 orders and almost by 4 orders while approaching the axis as compared to those at the periphery. The number of the formed layers is defined by the balance of energy and momentum.

All these considerations are concerned with nonrelativistic motions. But, in view of a high density of electrons in the layers arisen in the cases under study, the electron states must be described with relativistic equations.

Below, in the analysis of structures, we will use a generalization of the methods developed in Ref. 346 to the relativistic case. Instead of a kinetic equation for electrons, we use the Dirac equation for the wave function or relativistic hydrodynamic equations which will define the contribution of the density of electrons to the Poisson equation.

Thus, the Poisson equation becomes practically a modified Thomas–Fermi equation.

A Modified Thomas–Fermi Equation and Plasma-Field Structures.

With increase in the density of electrons, their distribution $n_e(r, \varphi(r))$ in the shell must be described by relativistic equations. The electrostatic potential inside and outside the shell can be found, as above, from the Poisson equation, where one should substitute the expressions for the density of electrons and nuclei in the shell in terms of the potential, which will reduce it to a nonlinear equation of the Thomas–Fermi type (see Ref. 349).

With regard to the relativistic coupling of energy and momentum (see Sec. 12.1.), the equations for the electron density and velocity in a quasistationary flow take the form

$$\operatorname{div} \left(n_e \frac{\vec{v}_e(\psi(r))}{c} \right) = 0, \quad \sqrt{x_e^2 + 1} - \psi = \omega_0, \quad (12.277)$$

where $\vec{v}_e/c = x_e(\psi)/\sqrt{1 + x_e(\psi)^2}$, $\psi = e\varphi/m_e c^2$, $x_e = p/m_e c$.

From here, we get the functional dependence between the density of electrons and the distribution of the potential in the shell:

$$n_e(r) = n_{e0} (v_{e0}/c) R_0^2 r^2 \sqrt{1 - 1/(\omega_0 + \psi(r))^2}. \quad (12.278)$$

The dependence of the density of nuclei on the potential is not changed as compared to Eq. 12.275, because the motion of nuclei remains nonrelativistic. Thus, we obtain a nonlinear equation for the potential which is a generalization of the Thomas–Fermi equation:

$$\frac{\partial}{\partial r} \left(r^2 \frac{\partial \psi}{\partial r} \right) = 4\pi \frac{e^2}{\hbar c} \left(\frac{\beta_e}{\sqrt{1 - 1/(\omega_0 + \psi(r))^2}} - \frac{\beta_i}{\sqrt{1 - \eta \psi(r)}} \right). \quad (12.279)$$

It is seen from Eq. 12.279 that the potential will be real if the values of the dimensionless potential vary between the following limits:

$$(1 - \omega_0) < \psi(r) < 1/\eta. \quad (12.280)$$

As distinct from a vacuum shell, where the spatial region with $r \rightarrow 0$ is most significant, the radii for all components of an electron-nucleus shell differ slightly from the shell radius R . In this case, it is convenient to introduce the new variable $\rho = R - r$. Because $\frac{\rho}{R} \ll 1$ almost always during the evolution of the shell, Eq. 12.279 is easily reduced to the plane case:

$$\frac{d}{d\rho} \left(\rho^2 \frac{d\psi}{d\rho} \right) = \frac{4\pi e^2}{R^2 \hbar c} \left(\frac{\beta_e}{\sqrt{1 - 1/(\omega_0 + \psi(\rho))^2}} - \frac{\beta_i}{\sqrt{1 - \eta \psi(\rho)}} \right). \quad (12.281)$$

By integrating this equation once, we get a nonlinear equation of the first order for the potential:

$$\left(\frac{\partial \psi}{\partial \rho} \right) = \pm \sqrt{\frac{2}{3\pi} \frac{e^2}{\hbar c} \left(\Psi_u(U) - \Psi_u(U|_{\rho=\rho_0}) \right) + \left((\partial U / \partial \rho)|_{\rho=\rho_0} \right)^2}, \quad (12.282)$$

where $\Psi_u(U) = \left(\beta_e \sqrt{(\omega + U)^2 - 1} + (2\beta_i/\eta) \sqrt{1 - \eta U} \right)$.

The solution of Eq. 12.281 can be written in the implicit form as

$$\int_{U(\rho)}^{U(\rho_0)} \frac{du}{\sqrt{\frac{2}{3\pi} \frac{e^2}{\hbar c} (\Psi_u(U(\rho_0)) - \Psi_u(u)) + \left((\partial U / \partial \rho)|_{\rho=\rho_0} \right)^2}} = \pm(\rho - \rho_0). \quad (12.283)$$

Similarly to the nonrelativistic case, this solution has character of an oscillatory process, whose different semiperiods are defined by different signs on the right-hand side.

The amplitude of oscillations of the potential in a structure formed in such a way increases, which is related to the system sphericity.

Plasma-Field Structures in the Nonequilibrium Electron-Nucleus Plasma. The appearance of electric fields with extreme intensity, as well as of the powerful flows of particles, in a target results in that the system under consideration is far from the equilibrium state. As a measure of remoteness of the system from the equilibrium, we may take the intensities of the field and the flows of particles and energy, which depend on the field intensity, in the phase space.

As indicated above, the Fermi–Dirac equilibrium distribution function should be modified in this case and can be approximated by the relation $\bar{p}_i = \frac{1}{Z_q} \left[\exp_q(-\beta(E_i - \mu Q_i)) \right]^q$, where the index of nonextensiveness q is defined by the electric field intensity in the formed structure. This modified distribution function has power “tails”.

The imbalance of the system leads to a more complicated pattern of the evolution of the plasma-field structure, because one more parameter arises, but without any qualitative changes in this pattern.

We recall that, due to the spherical symmetry, the potential in the structure increases sharply while approaching the center. Therefore, for the next layer positioned closer to the center, the energy of electrons, being able to penetrate into the layer, should be sufficiently high. This yields that the structure size is defined by the length of the “tail” of the distribution function in the energy space. It is obvious that a significant increase in the number of electrons in the tail of the distribution function results in an increase in the number of layers in the structure. For a power asymptotics of the distribution function, the energy increases by orders of magnitude.

Since the parameter q is defined by quantities which can be controlled during experiments, it is clear that the described process allows one to efficiently control the formed structure.

As a result, there arises a lot of layers with characteristic spatial size $\sim 10^{-12}$ cm and with the densities of nuclei and electrons $\sim 10^{32}$ cm $^{-3}$ and $\sim 10^{34}$ cm $^{-3}$, respectively. In the system of the formed layers with a high density of nuclei, the mean distance between nucleons is about 30 to 100 fm. For these parameters of the layers, the Coulomb field of nuclei is compensated at the expense of the Debye screening by electrons of the medium.

12.3. Binding Energy of Nuclear Systems and Nonequilibrium Thermodynamics

The investigation carried out in the previous Sec. 12.2 has shown that the evolution of the dynamical and electrodynamical processes running in a target under the action of an external pulse REB leads, on the final nonlinear stage of its development, to a specific two-fluid mode of the flow of a substance. This mode is the most bright manifestation of collective properties of the electron-nucleus system and the end of the passive part of the evolution of this self-consistent nonlinear wave of the volume charge density.

The arisen mode of the flow of a substance is the propagation of a nonlinear electron Langmuir wave along the ion flow. In intense electron flows, the wave steepens rapidly (see Sec. 12.2.4), which is accompanied by the appearance of great local densities of electrons and electric fields on its leading edge (see Sec. 12.2.4).

In its final stage, the evolution of a plasma wave leads to the formation of plasma-field structures with small spatial scales and great fields (see Fig. 12.62). The evolution of a plasma-field structure consists in the collapse of nonlinear Langmuir waves into the system of superthin concentric layers, which propagates toward the target center.

The ordinary collapse of Langmuir waves represents both the process of contraction of all plasmons to the geometric center of the system and the simultaneous extrusion of a substance from the central region of the system by the electromagnetic field.

In our case, the Langmuir collapse possesses a number of unique properties. *The collapse of plasmons yields a multidimensional structure rather than a 1D one.* The surface, to which plasmons are moving (the region with an extremely intense electric field), is surrounded by surfaces, to which the particles with opposite charges (electrons and nuclei) are moving. As shown in the previous section, these surfaces are a fractal structure with a spatial scale reducing toward the center.

The arisen plasma-field structures are the series of numerous thin spherical electronic and nuclear layers embedded one into another. The density of nuclei of about $10^{29}/Z$ cm $^{-3}$ is concentrated in narrow layers with

thicknesses $\delta_{str} \approx 1.5 \cdot 10^{-10}$ to $2.0 \cdot 10^{-10}$ cm. The layers of electrons have thicknesses greater than those of the layers of nuclei.

The movement of the structure toward the center ensures the possibility to increase the substance density. We emphasize the peculiarities of the evolution of the above-described Langmuir collapse by using, in what follows, the term “a two-cascade Langmuir collapse” for it.

As shown in Part II devoted to the analysis of experimental data, we may say about the appearance of a high-density high-temperature region with micron size in the central part of a target as a result of the evolution. For brevity, we will call this region as a “hot dot”. A hot dot is the point-like source of X-ray emission, whose typical spectrum is shown in Fig. 12.63.

The registration of the X-ray emission spectrum from the active region of the anode allows us to estimate the important characteristics of a nonequilibrium state of HD.

The numerical quantum-mechanical calculations of the bremsstrahlung for strongly screened potentials (Thomas–Fermi and more complex ones) showed that the *X-ray emission intensity is suppressed at low frequencies*. The molecular-dynamical calculations of the effect of interion correlation also demonstrated a weakening of the emission at low frequencies. It is clear because *low frequencies correspond to long paths, which are not realized practically in dense media*.

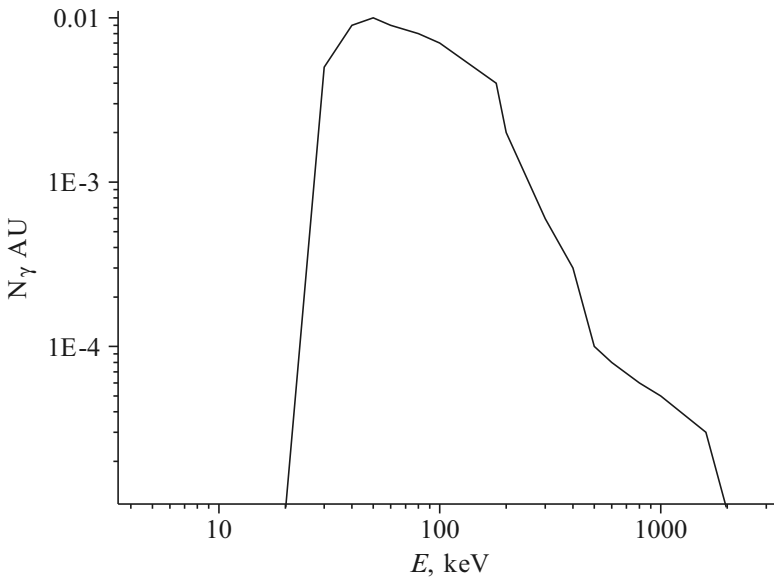


Fig. 12.63. Experimental spectrum of the X-ray emission from HD on the double logarithmic scale.

The position of the X-ray spectrum maximum allows us to estimate the temperature of the surface of a dense hot region as $T_i = 38 \text{ keV}$.

We can also indicate the characteristic frequency ω_{cr} , below which the continuum is suppressed. This limit frequency $\omega_{cr} = \omega^* \gamma_{ei}^{3/2}$, where $\omega^* = (6\pi e^2 Z n_i / m)^{1/2}$ and $\gamma_{ei} = Ze^2 / \bar{r} T$ is a parameter of the electron-ion interaction which is mainly defined by the plasma frequency of a plasma, being present in HD.

The expression for the limit frequency and the above-presented estimate of the temperature of HD located in the anode allow us to estimate the density of this region. Indeed, the experimental studies of the X-ray emission spectrum show that the spectrum is sharply cut off near the energies of quanta of about 1 keV. By using the estimate $\omega_{cr} = 1 \text{ keV}$ and Eq. 12.286, we get from Eq. 12.285 that $n_i = 1.73 \cdot 10^{27} \text{ cm}^{-3}$. The presented estimates are the lower bounds for the temperature and density of HD.

Large values of the density in structures arising upon a two-cascade Langmuir collapse and peculiarities of these plasma-field structures lead to the essential influence of the effects of polarization and screening on the interaction of nuclei. The corresponding changes are so significant that the character of nuclear reactions sharply changes, and the latter become properly collective.

Below, we show that the size effect arising due to a small thickness of shells in the structure leads to the effective coupling of nuclei in its volume and to the union of all nuclei from concentric shells into a structure which is, in the first approximation, a multilayer spherical fullerene-like cluster consisting of α -particles.

A significant change in the rate of the reactions of synthesis is possible only upon the joint action of several factors: the increase of the density, degree of imbalance, and intensity of collective self-consistent fields. Under these conditions, the probability of the reaction of synthesis increases sharply and approaches unity at moderate temperatures (see Sec. 1.3).

It is worth noting that the theoretical investigations, which are carrying on at the Electrodynamics Laboratory "Proton-21", of nuclear processes agree with the plan of the development of nuclear physics in Europe. This plan was developed at the meeting of famous European scientists intensively working in the field of nuclear physics. In the summary of this meeting, a great attention was given to various nuclear structures (to both quite traditional and sufficiently exotic ones). It was also noted that many assertions which seemed earlier to be true forever are not valid at present or are violated in some cases.

As such an assertion, we mention the well-known relation $R = r_0 A^{1/3}$ for the radius of a nucleus R as a function of the mass number A . The very

general geometric consideration allows one to conclude that this relation holds only under the simple assumption about the homogeneity of a distribution of the nuclear substance in a nucleus. The possibility of a strong inhomogeneity of the nuclear substance even in ordinary nuclei was assumed long ago (see, e.g., Refs. 350, 351). The calculations executed on the basis of the Hartree–Fock method with a velocity-dependent potential allow one to suppose that nuclei ^{36}Ar , ^{200}Hg , and others can turn out to be bubble-like.

Here, we try to clarify some quantum and macroscopic properties of nuclear clusters which can be most important for the realization of nucleosynthesis. To this end, we consider

- peculiarities of the description of multiparticle, kinetic, and hydrodynamic processes in nuclear systems
- peculiarities of the polarization-involved interactions and their influence on the rates of the reactions of synthesis of nuclei and the growth of a nuclear cluster
- estimates of the energy of nuclear clusters with regard to strong correlations in them
- effects of condensation of fermions and bosons and their influence on the energy of nuclear clusters
- main relations of the nonequilibrium thermodynamics of nuclear clusters and the processes of evaporation of nuclei

The description of the dynamics of a nuclear system under nonequilibrium conditions is a complicated problem which concerns the internal properties of the system, its energy spectrum (the ground state energy and the spectrum of excited states), mean values of dynamical variables, their distributions, probabilities of transitions, etc.

In the period of its formation, quantum mechanics was considered as the physics of the microworld, whereas macroscopic phenomena remain to be described by classical mechanics. For recent decades, we observe a growth of interest in macroscopic phenomena which are directly related to quantum mechanics. They are, first of all, the phenomena connected to the notions of collective variables and coherency.

An important problem is the description of dissipative and open quantum systems. The achievements attained in the description of classical open systems have led to the creation of such a field of scientific knowledge as synergetics.

At present, very intense are the studies of the processes of self-organization in quantum systems under dissipation and upon the simultaneous action of external forces on a system. Most important for the description of the real world are nonequilibrium and nonstationary quantum systems.

An increase in the degree of imbalance in an open system is related to a decrease in its entropy and the appearance of correlations and structures. In the course of a nonequilibrium process, a spatial structure can spontaneously arise from a spatially homogeneous state, i.e., we are faced with the phenomenon of self-organization.

The approach to open systems on the base of quantum theory requires from the very beginning to introduce new objects such as a density matrix, statistical operators, and Green functions which transfer the consideration, in fact, in the field of quantum statistics.

Upon the self-organization, some order appears in a dissipative system, and the dynamics acquires the collective character, which gives the possibility to intensify the flows of various quantities at the expense of an ordered macroscopic motion. In a quantum dissipative system, the phasing and coherent phenomena play a significant role.

Dissipative structures appear only in the systems which are described by nonlinear equations for macroscopic variables. Upon the study of such systems, we meet a complicated interlacing of the statistical and informational aspects of quantum theory.

At present, the description of similar situations involves generalizations of nonrelativistic quantum mechanics. In this case, the principle of correspondence yields that the Schrödinger equation corresponds to the Liouville equation, and nonlinear equations of quantum mechanics of the Hartree–Fock type do to kinetic equations (e.g., the equations of Vlasov, Boltzmann, and others) (see Ref. 323). A very interesting general theory for the description of the thermodynamical properties of systems (especially for the systems near phase transitions) was constructed by V.P. Maslov in Ref. 383. He demonstrated, in particular, the essential role of the discreteness of a system of particles and the efficiency of the representation of the system of particles in the form of enumerated pairs, which leads to the Schrödinger equation, as well as to the kinetic equations for the Wigner function. This result allows one to assume that the attempt to establish the direct connection between quantum statistical theory and the control theory for discrete multiconnected dynamical systems will be perspective.

12.3.1. Kinetic and Hydrodynamic Equations for the Nuclear Matter: Nonequilibrium Stationary States of Nuclear Particles

The basis of the modern description of the systems of many particles in the quantum region is given by the field methods of the quantum theory of many particles. Generally, these methods are based on the ordinary N -particle Schrödinger equation

$$\left(i\hbar \frac{\partial}{\partial t} - H\right) \psi(\mathbf{q}, t) = 0. \quad (12.284)$$

Here, $\mathbf{q} = \{q_1, \dots, q_N\}$ is the collection of coordinates of the system of N particles; q_i stand for the spatial coordinates x_i and the discrete (the spin coordinates σ_i and the isotopic ones τ_i) ones of the i -th particle. The Hamiltonian H is considered to be a given function of the coordinates, operators of momentum $\mathbf{p} = -i\nabla$, and the matrices that transform the discrete coordinates.

The Hamiltonian of the system can be represented as the sum of two terms: $H = H_F + H_I$. In the simplest case, the free Hamiltonian H_F is the sum of the operators of the kinetic energy of particles and those of the energy of their interaction with an external field U :

$$H_F = \sum_{i=1}^N T_i, \quad T_i = \frac{p^2}{2m} + U(q); \quad (12.285)$$

the interaction potential is

$$H_I = \frac{1}{2} \sum_{i < j}^N V_{ij}, \quad V_{ij} = \frac{Z^2 e^2}{r_{ij}}, \quad r_{ij} = |\mathbf{x}_i - \mathbf{x}_j|, \quad (12.286)$$

where Ze is the nucleus charge.

The interaction potential between nuclei and nucleons depends also on the operators of momentum, spin, etc., and, hence, the very character of the interaction between particles (its intensity and sign) depends on their state. One may expect the appearance of multiparticle forces (three-body, four-body, etc.) between nucleons.

The methods of quantum field theory are the direct and natural development of the method of secondary quantization. Of the primary importance become the operators of annihilation and creation of particles $a(x)$ and $a^+(x)$. The process of interaction between particles can be interpreted as the annihilation of particles in the initial state and their creation in the final state.

The natural tool for the study of systems of many interacting particles is the Liouville–von Neumann equation for the density matrix following from the Schrödinger equation Eq. 12.284:

$$i\hbar \frac{\partial}{\partial t} \rho(t) = [H(t), \rho(t)]. \quad (12.287)$$

The Wigner transformation

$$n(p, r) = \int \frac{d^3 k}{(2\pi\hbar)^3} \exp(ikr) a_{p-k/2}^+ a_{p+k/2} \quad (12.288)$$

leads after the averaging, to the Wigner function $f^W(p, r, t) = \langle n(p, r) \rangle^t$ in phase space (see Refs. 277, 353). The Wigner function is suitable, because it is simply connected with the classical distribution function through the averaging over the elementary cell of phase space with volume $\hbar^3$:

$$\frac{1}{\Delta\mu} \int f^W(p, r, t) d\mu = f(p, r, t) + O\left(\frac{\hbar^3}{\Delta\mu}\right). \quad (12.289)$$

By using these definitions, we can reduce the equation for the Liouville density matrix Eq. 12.287 in the case of binary collisions to the form coinciding with that of the kinetic equation for a one-particle distribution function

$$\partial f(\mathbf{p}, \mathbf{r}, t) / \partial t = I_{st}(\mathbf{p}, \mathbf{r}, t), \quad (12.290)$$

where the collision integral reads

$$\begin{aligned} I_{st}(\mathbf{p}, \mathbf{r}, t) = & \frac{2\pi}{\hbar} \int \frac{d^3 p_2 d^3 p'_1 d^3 p'_2}{(2\pi\hbar)^3} W_{p_1 p_2 \rightarrow p'_1 p'_2} \\ & \times \delta(E_p + E_{p_2} - E_{p'_1} - E_{p'_2}) \delta(\mathbf{p} + \mathbf{p}_2 - \mathbf{p}'_1 - \mathbf{p}'_2) \\ & \times \left\{ f_p f_{p_2} (1 - f_{p'_1}) (1 - f_{p'_2}) - f_{p'_1} f_{p'_2} (1 - f_p) (1 - f_{p_2}) \right\}. \end{aligned} \quad (12.291)$$

When the density and correlations increase in a system of nuclear particles, it is necessary to consider multiparticle effects and the formation of bound states, i.e., the states with long-lived correlations of particles in the system.

From the macroscopic viewpoint, the existence of stable bound states in the system can be described as the formation of particles of a new type, being the clusters formed from initial particles. In such an approach, it is possible to study not only the scattering of particles, but also the reactions running in the system of initial particles and formed clusters.

Further, for the kinetic description of our system, we will use the simplest estimates and get the main characteristics of nonequilibrium states of the nuclear system by separating the evolution in the energy space from the evolution in the configuration space. First of all, we will obtain the quasistationary solutions of equations for the one-particle distribution function Eq. 12.290 with the collision integral Eq. 12.291.

Nonequilibrium Stationary States of Nuclear Particles. As usual, the state of a system moves away from the equilibrium one with increase in the flows of energy or particles present in the system. Well studied are the locally nonequilibrium states. In these states in every physically arbitrary small volume separated around any point of space, there exists an equilibrium state, whose parameters vary from point to point. The gradients of

these parameters define both the flows of various quantities in space and the degree of a deviation from the equilibrium.

A more complicated situation arises in spatially homogeneous systems or in those regions of the system which can be approximately considered homogeneous. In such regions of the system, a deviation from the equilibrium is already related to the impossibility to use the approximation of local equilibrium rather than to a change of the equilibrium parameters in space. *In this case, the deviation from the equilibrium is defined by the presence of flows not in the configuration space, but in the full phase one* (see Refs. 314, 342). *Intensities of these flows characterize the degree of deviation of a quasistationary nonequilibrium state of the system from its equilibrium state.*

The importance of nonequilibrium statistical effects in nuclear physics is related to the extremely high density of nuclei, a growth of the memory effects at such densities, and long-range forces. These effects imply the presence of non-Markov processes in the kinetic equation describing the process of transition to quasistationary states and the very quasistationary states.

The rigorous definition of the states with nonextensive properties, which arise due to the memory effects and/or long-range interactions, should be based on the microscopic equations of a nucleon-nucleus plasma such, for example, as a kinetic equation for the one-particle distribution function. Consider this equation in the spatially homogeneous case. The collision integral in the Boltzmann form can be rewritten in the Landau form for the Coulomb interaction (see Ref. 357), i.e., in the form of the derivative of the flow of particles in phase space $J_{nl}\{f, v\}$:

$$I_{nlFP}[f, f] = -\frac{1}{4\pi v^2} \frac{\partial}{\partial v} J_{nl}. \quad (12.292)$$

For the flow of particles in the phase space J_{nl} , we can get different approximations, e.g., $J_{nl} = \Pi_{nl}(v, \{f(v)\})$. Most convenient for the analysis is the expression for the flow reduced to a symmetric form. In classical statistics, we have (see Ref. 314)

$$\begin{aligned} \Pi_{nl}(v, \{f(v)\}) &= -\frac{4\pi}{v} \frac{\partial}{\partial v} \int_0^v f(v)P(x) - f(x)P(v)]x^2 dx, \\ P(v) &= 2 \int_v^\infty f(x)xdx. \end{aligned} \quad (12.293)$$

In quantum statistics, Eq. 12.292 is easily generalized by the introduction of the Pauli multipliers. In this case, the kinetic equation becomes nonlinear, even if the integral coefficients in Eq. 12.293 are linearized and turn out to be independent of the distribution function. The kinetic equation acquires the form of a Fokker-Planck nonlinear differential equation.

At the great intensities of sources, it is necessary to consider the processes of relaxation of the flow between collisions. The relaxation of the flow leads to the necessity to pass from the standard parabolic Landau equation to the system of equations of the hyperbolic type:

$$\frac{\partial f}{\partial t} = I_{nlFP}[f, f] + \Psi(v); \quad \tau \frac{\partial J_{nl}}{\partial t} + J_{nl} = \Pi_{nl}(v, \{f(v)\}). \quad (12.294)$$

For the first time, the kinetic equation in the differential form in the frame of quantum statistics was deduced by A. Kompaneets (Ref. 358) for a photon gas interacting with an electron gas; it reads

$$\frac{\partial f}{\partial \tau} = \frac{\nu_0 h}{mc^2} \frac{1}{\nu^2} \frac{\partial}{\partial \nu} \nu^4 \left(\frac{T}{h} \frac{\partial f}{\partial \nu} + f(1+f) \right), \quad \nu_0 = \sigma_T n_e c. \quad (12.295)$$

Upon the derivation of this equation, only the statistical properties of the system were directly used. The character of the interaction was manifested only in the dependence of the frequency of collisions on the energy. This circumstance allows one to apply an equation of this type not only to the scattering of photons by electrons, but to any system of bosons undergoing a collision with a system of heavier particles.

It was shown in our works (see, e.g., Refs. 314 and 342) that the Fermi–Dirac equilibrium distribution function can be generalized for states with strong correlations to read

$$f_q(\varepsilon, E_f, \beta) = \frac{1}{1 + \exp_q(\beta(\varepsilon - E_f))} \theta(E_f - \varepsilon) + \frac{\exp_q(-\beta(\varepsilon - E_f))}{\exp_q(-\beta(\varepsilon - E_f)) + 1} \theta(\varepsilon - E_f). \quad (12.296)$$

Figure 12.64 demonstrates the behavior of the distribution function $f_q(\varepsilon, E_f, \beta)$ as a function of energy upon a change of the parameter of nonextensiveness q .

This distribution function is a stationary solution of the kinetic equation, which is seen from the analysis of Fig. 12.65. In this figure, we draw the surface in the space (v, q) which represents a flow in the phase space $\Pi_{nl}(\varepsilon, \{f_q(\varepsilon, E_f, \beta)\})$ with regard to quantum corrections calculated with the use of the distribution function Eq. 12.296.

Stationary states correspond to the regions, in which the flow does not depend on the velocity. It is seen that there are two such regions.

The first region with $q = 1$ corresponds to a Maxwell's equilibrium solution, whereas the second one does to the states with power asymptotics, which were first derived by us in Ref. 313.

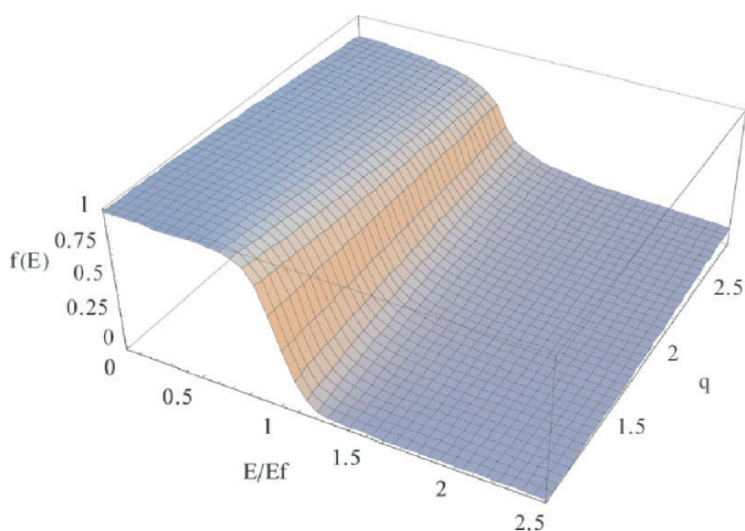


Fig. 12.64. Dependence of the distribution function on energy and the parameter of nonextensiveness q .

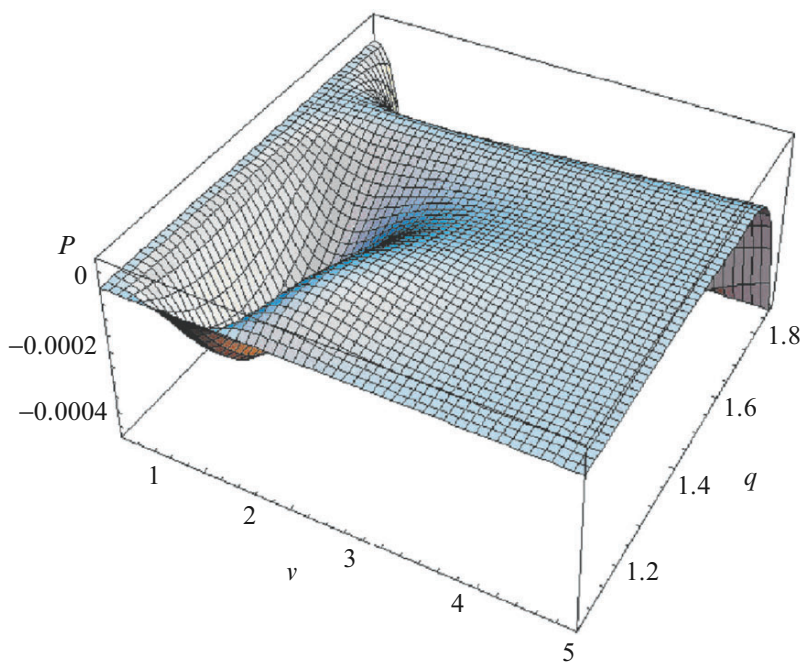


Fig. 12.65. Dependence of the flow of particles in the phase space on the velocity and the parameter of nonextensiveness q .

Possessing the distribution function, we can find the density of particles

$$Q(E_f, \beta, q) = \int_0^\infty g(\varepsilon) f_q(\varepsilon, E_f, \beta) d\varepsilon \quad (12.297)$$

and the energies

$$E(E_f, \beta, q) = \int_0^\infty g(\varepsilon) \varepsilon f_q(\varepsilon, E_f, \beta) d\varepsilon, \quad (12.298)$$

where $g(\varepsilon)$ is the density of states of particles in the system.

The dependence of the system energy on the parameter q is shown in Fig. 12.66(a).

As usual, the parameter β in the distribution is simply expressed via the physical temperature in the system of particles ($\beta = 1/T$). For the nonequilibrium case, it was shown that a more complicated relation between β and the physical temperature $\bar{T}$ is valid:

$$\bar{T} = \beta^{-1} + (q - 1)(E - \mu Q). \quad (12.299)$$

It is presented in Fig. 12.66(b).

Above, we have considered the peculiarities of the nonequilibrium quasistationary distributions of fermions. As for the peculiarities of the distribution functions of bosons under nonequilibrium conditions, it is convenient to analyze them in the framework of the differential kinetic equation Eq. 12.295 which was deduced and studied by A. Kompaneets for photons without any sources.

For nonzero values of the flow defined by the integral $\int_0^\varepsilon d\varepsilon' \Psi(\varepsilon')$, the stationary nonequilibrium distribution function of photons can be defined, as shown by us (see, e.g., Ref. 360), from the equation

$$\nu(\varepsilon)g(\varepsilon) \left[T_e \frac{\partial f}{\partial \varepsilon} + f + \theta f^2 \right] = \int_0^\varepsilon d\varepsilon' \Psi(\varepsilon'). \quad (12.300)$$

Here T_e is the temperature of electrons. The solution of Eq. 12.300 corresponds to a nonequilibrium quasistationary state with a constant flow of the number of particles in the energy space.

We note that, from our viewpoint, it is possible to use the electromagnetic analogies (see Ref. 361) for the description of the kinetics of pions upon their interaction with nucleons due to the smallness of the pion mass

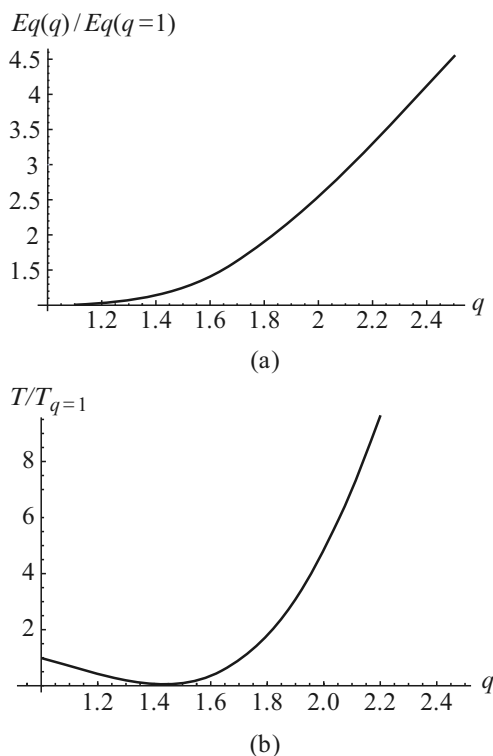


Fig. 12.66. Dependence of the ratio of the energy density in the system of particles to the equilibrium density on the parameter q (a); dependence of the ratio of the physical temperature to the temperature in the equilibrium on the parameter of nonextensiveness q (the parameter of correlation) in the system of particles (b).

as compared to the masses of nucleons. So, an equation similar to the Kompaneets equation should be true for pions. Similarly to the description of the kinetics of the Bose-condensation of photons upon their interaction with electrons by the Kompaneets equation, the given modified equation should describe the kinetics of the processes of condensation of pions upon their interaction with nucleons.

As follows from the analysis of works on the study of shock waves in the spectrum of photons (see Ref. 358), we may expect the formation of a fine structure of the leading edge of a shock wave upon the Bose-condensation of pions as well.

The Uncertainty Relations for Coherent States and States with Strong Correlations. The analysis of consequences of the uncertainty relations for the states, which are coherent along the radius and possess

strong correlations on the surface of a shell, was carried out in the general case in Chapter 2. The knowledge of the above-presented correlation characteristics of states in a shell nuclear cluster allows us to sharpen the relevant estimates.

As seen from the analysis of the results of our investigation, the states of particles in a shell possess the basically different properties in the radial direction and on the shell surface (these properties have been efficiently used in Chapter 2). Such a difference in the properties along the radius and on the shell surface leads also to the necessity to rewrite the uncertainty relation in the essentially different forms for different coordinates of the physical space. Since the states of particles have a strong correlation on the surface, the uncertainty relation on the surface should be modified with regard to correlations (see Ref. 21) and the mixing of states.

Correlations in the system of particles are defined by the coefficient of correlations $C_c^2 = \frac{\langle(\Delta p \Delta x)^2\rangle}{\langle(\Delta p)^2\rangle\langle(\Delta x)^2\rangle}$ varying in the interval $0 \leq C_c^2 < 1$, and the mixing of states can be characterized by different quantities. For example, it is convenient to characterize the purity of states by the entropy $S = -\int f \ln(f) d\Omega$ (here, f is the distribution function in the phase space, and $d\Omega$ is an element of the phase space). In this case, we have

$$\begin{aligned}\sqrt{1 - C_c^2} [\Delta\psi \cdot r(t)] \cdot \Delta p_\theta(t) &\geq \frac{\hbar}{2} F(S), \\ \sqrt{1 - C_c^2} [\Delta\varphi \cdot r(t)] \cdot \Delta p_\varphi(t) &\geq \frac{\hbar}{2} F(S),\end{aligned}\quad (12.301)$$

$$F(S) = \left(1 + \frac{2}{\exp(\beta(S)) - 1}\right), \quad \frac{\beta(S)}{\exp(\beta(S)) - 1} - \ln(1 - \exp(-\beta(S))) = S.$$

For the states with strong correlations, $S \gg 1$, and the function on the right-hand side of the inequality can be simplified: $F(S) = 2 \exp(S - 1)$.

In view of the inequalities $\Delta\psi \cdot r(t) \leq 2\pi r(t)$ and $\Delta\varphi \cdot r(t) \leq 2\pi r(t)$, relations Eqs. 12.301 yield

$$\Delta p_{\varphi, \Psi}(t) \geq \frac{\hbar}{4\pi r(t) \sqrt{1 - C_c^2}} \exp(S - 1). \quad (12.302)$$

Here, the entropy S should be replaced by its nonextensive generalization S_q . The parameter of nonextensiveness q defines the correlations in the system and will be estimated below from the comparison with experimental data. The appearance of a structure is accompanied by a growth of correlations and the quantity S_q . As seen from the structure of the uncertainty relation Eq. 12.302, a positive feedback appears between a decrease in the shell radius and a growth of correlations (a growth of S_q). Such a

feedback testifies to the development of a phase transition and the explosive growth of the dispersion of angular coordinates $\Delta\psi$ and $\Delta\phi$. That is, we may expect the realization of a fast explosive “smearing” of the region of the overlapping of wave functions of all the strongly correlated particles over the 2D surface of the shell. The overlapping of wave functions leads to a strong interaction of nucleons. It is difficult to say definitely about the effect of such strong correlation on the main types of the interaction of nucleons, but we consider below how the electromagnetic interactions ensure the indicated phase transition related to a modification of the Coulomb interaction of nuclei as a result of the exchange interaction in a thin plasma shell.

In a 1D subspace $r(t)$, the state of particles of the ensemble is coherent, and the uncertainty relation for coordinates–momenta degenerates, as known, in the equality

$$\Delta r(t) \cdot \Delta p_r(t) = \hbar/2. \quad (12.303)$$

This relation yields that, as long as the shell structure exists ($\Delta r(t) \ll r(t)$), the momentum dispersion $\Delta p_r(t)$ increases with decrease in $r(t)$ and shows a singularity as $r(t) \rightarrow 0$. This change in the momentum dispersion is inevitably accompanied by a sharp increase in the kinetic energy of the motion of a plasma-field shell structure, which is directed toward the shell center, and corresponds to its collapse. We now pass to the analysis of the physical mechanism of a phase transition in thin plasma concentric layers.

12.3.2. Influence of Dynamical Polarization on Pycnonuclear Reactions and the Growth of Clusters

Collective effects play a significant role upon the interaction of nuclei and nucleons and lead to the interesting peculiarities which should be taken into account upon the description of the processes of clusterization. The influence of the static screening and polarization on the interaction of nuclei and the rate of reactions was studied in many works (see, e.g., Refs. 362, 363).

In the simple model of a spherical nucleus containing A nucleons (Z protons and $N = A - Z$ neutrons), the normal substance density is of the order of $\rho_0 = 0.14 \text{ fm}^{-3}$. This density corresponds to the radius of a spherical nucleus

$$R_A = \left(\frac{A}{\frac{4}{3}\pi\rho_0} \right)^{1/3} = 1.2 A^{1/3} \text{ fm}. \quad (12.304)$$

In this case, if the distance between nuclei significantly exceeds their radii, they interact mainly via the repulsion force. When nuclei (charged drops) approach one another at a distance which is less or of the order of the radius of screening of the Coulomb field of nuclei, there appear the

redistribution of charges in nuclei and the polarization of nuclei. As known (Ref. 364), the polarization is described by the operator of dipole moment of a nucleus $\mathbf{d} = e \sum \mathbf{r}_p$, where the summation is performed over all the protons in a nucleus. The dipole moment can be written in more symmetric form as

$$\mathbf{d} = e \left(1 - \frac{Z}{A} \right) \sum_p \mathbf{r}_p - e \frac{Z}{A} \sum_n \mathbf{r}_n, \quad (12.305)$$

with the summation over protons and neutrons. Protons and neutrons have, respectively, the “effective charges” $e(1 - Z/A)$ and $-eZ/A$.

Consider two nuclei approaching each other. The charge of the first nucleus creates the electrostatic image in the second one, which becomes, in its turn, the source of an image in the first nucleus (see Ref. 362). Thus, the field intensity at an arbitrary point between nuclei on the axis joining their centers is created by the infinite number of the images formed in both nuclei. The polarization forces change the sign of the Coulomb interaction between nuclei upon their approaching.

The description of the processes of binary collisions of nuclei in the frame of quantum mechanics includes the tunneling of particles through the Coulomb barrier, which is defined by the interaction potential with regard to the contribution of the polarization interaction, being proportional to the dipole moment:

$$U(r) = \frac{Z_1 Z_2 e^2}{r} - \frac{d_z Z_2 e}{r^2}. \quad (12.306)$$

We note that the Schrödinger equation for this potential was considered in Ref. 364.

A nucleus polarized in the external electric field has dipole moment. The polarization of nuclei in external fields makes the Coulomb barrier thinner and increases exponentially the probability of the process of tunneling. The tunneling probability increases in the high-density region and in collective fields. However, changes in the Coulomb interaction of charged particles turn out, as before, only corrections which induce the corrections in the probabilities of the tunneling and in the rates of nuclear reactions. These corrections are revealed, despite their smallness, in sufficiently great time intervals and are essential upon the study of the proton–nucleus reactions in astrophysics.

Reference 366 contains the analysis of the screening of the Coulomb interaction in an unbounded equilibrium plasma. There it was shown that, contrary to the results in Refs. 271 and 272, just the dynamical screening is most essential in a dense plasma rather than the static one.

The dynamical screening describes the influence of the polarization of a medium with regard to the excitation of its natural oscillations on the

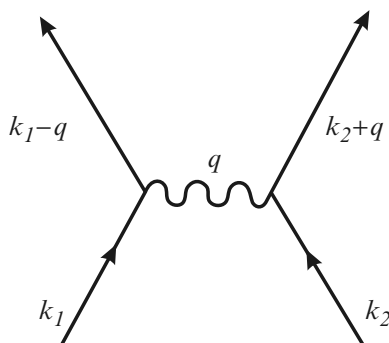


Fig. 12.67. Diagram of the interaction of nuclei via plasmons.

interaction of point-like charges. In the quantum language, there arises the exchange interaction of charged particles via virtual quanta corresponding to natural oscillations of the medium. The simplest diagram describing such an interaction is given in Fig. 12.67.

A nucleus with momentum k_1 emits a quantum with momentum q and remains with momentum $k_1 - q$, and the second nucleus with momentum k_2 absorbs this quantum with momentum q at other point of space and at the other time moment, and its momentum becomes $k_2 + q$. There appears the attraction between these nuclei.

Under certain conditions, these collective interactions induce a phase transition where there occurs the qualitative change of the interaction rather than a small one. For example, just the modification of the Coulomb interaction of electrons in the plasma of a solid as a result of the dynamical polarization-involved processes causes the attraction between electrons, which leads eventually to their pairing and to the phenomenon of superconductivity at temperatures below the critical one. These processes are most simply described in the quasiclassical “plasma” language (see Ref. 365).

We get a successive complete pattern of the Coulomb interactions of charges with regard to collective interactions in the medium surrounding charges, if the static screening (the ordinary Debye screening) and the dynamical screening of the Coulomb interaction in a dense plasma are described with the help of the Lennard–Balescu collision integral (see Refs. 341, 284, 366).

It follows from this approach that all particles in plasma can be considered as dynamically screened plasma excitations. While moving, these particles are transporting their own polarization “coats”. In the statistical equilibrium, the collision integral is turned to zero by the distributions which describe these plasma particles together with their polarization “coats”. Only a screened ion is a true plasma particle.

As mentioned above (see Sec. 12.2), the distributions, being solutions of a kinetic equation with regard to both the flows of particles in the phase space and correlations, are q -distributions in the form Eq. 12.297 for electrons and in the form Eq. 12.163 for nuclei. Namely they will be used by us upon the study of the rates of reactions. The interaction of particles is accounted for by the permittivity (see Refs. 341 and 365)

$$\varepsilon(k, \omega) = 1 + \frac{4\pi e^2}{k^2} \left[\int \frac{d\mathbf{p}}{(2\pi)^3} \frac{1}{\omega - \mathbf{k}\mathbf{v} + i0} \left(\mathbf{k} \frac{\partial f_e}{\partial \mathbf{p}} \right) + Z^2 \int \frac{d\mathbf{p}}{(2\pi)^3} \frac{1}{\omega - \mathbf{k}\mathbf{v} + i0} \left(\mathbf{k} \frac{\partial f_i}{\partial \mathbf{p}} \right) \right]. \quad (12.307)$$

We will illustrate the effect of screening by the expression for the Fourier transform of the potential $\varphi_{\mathbf{k},\omega}(\mathbf{v})$ of a test particle moving with velocity $\mathbf{v}$ (see Ref. 366) with regard to all polarization-involved processes defining the permittivity $\varepsilon(\mathbf{k}, \omega)$:

$$\varphi_{\mathbf{k},\omega}(\mathbf{v}) = \frac{Ze}{2\pi^2 k^2 \varepsilon(\mathbf{k}, \omega)} \delta(\omega - \mathbf{k}\mathbf{v}). \quad (12.308)$$

The corresponding potential of the Coulomb field, being a function of the coordinates and time, is given by the formula

$$\varphi(\mathbf{r}, \mathbf{v}, t) = \frac{Ze}{2\pi^2} \int \frac{d\mathbf{k}}{k^2 \varepsilon(\mathbf{k}, \omega)|_{\omega=\mathbf{k}\mathbf{v}}} \exp(i\mathbf{k}(\mathbf{r} - \mathbf{v}t)). \quad (12.309)$$

Upon the study of the processes of interaction of charges, the relations for an equilibrium unbounded plasma are usually used.

In our case, the dense plasma is, first, nonequilibrium and, secondly, bounded and is a system of moving thin plasma layers. The dispersion relations for thin plasma layers possess important peculiarities discussed in Refs. 367, 368.

Acoustic Plasma Waves in Thin Films: Film Quantization. As known, an equilibrium plasma consisting of at least two groups of charged particles allows the weakly decaying acoustic plasma waves to propagate (see Refs. 368–371). Physically, these waves correspond to the oscillations of the density of particles from separate groups with phase shifts providing the quasineutrality under summary oscillations.

In a homogeneous unbounded plasma, the damping of plasma waves can be small, if all particles of the plasma can be divided into at least two groups which have significantly different velocities in the propagation direction of a wave (see Refs. 368–370). Such a division can appear in the

system by various physical reasons. Acoustic plasmons with small damping, in which the presence of many components differing in energy results from the existence of the flows of particles in the phase space, were considered in Ref. 370. An analogous division into the groups of particles occurs in thin shells (films).

The transverse motion of the current carriers in shells (the motion along a radius) is quantized. In this case, the energy of a particle is defined by both the quasimomentum lying in the film plane and some discrete quantum number characterizing the energy of the motion being normal to the shell surface. Hence, the current carriers are divided into groups differing by a quantum number and the characteristic velocities along the film. In such a system, we can expect the appearance of acoustic plasma waves.

Consider the plasma waves propagating along the film under the film quantization (see Ref. 371). We assume that the electron gas in the film is strongly degenerated. In sufficiently thin films, there exist several branches of weakly decaying plasma waves with acoustic spectrum.

If we neglect an inhomogeneity near boundaries, the longitudinal permittivity ε^l is defined by the formula (see Ref. 368)

$$\varepsilon^l(q, q_s, \omega) = 1 + \frac{8\pi e^2}{(q^2 + q_s^2)} \sum_n \int \frac{d^2k}{(2\pi)^2} \frac{f(\varepsilon_{n+\nu, k+q}) - f(\varepsilon_{n, k})}{\hbar\omega + \varepsilon_{n, k} - \varepsilon_{n+\nu, k+q} + i\delta}, \quad (12.310)$$

where $\varepsilon_{n, k} = \varepsilon_n + \hbar^2 k^2 / 2m_{||}$ is the electron energy in the film, q is the wave vector of a plasmon lying on the shell surface, ε_n is the energy of a quantum transverse motion of an electron (for estimates, we take $\varepsilon_n = (\hbar^2 / 2m_{\perp}) (\pi n / \Delta_s)^2$), $n = \pm 1, \pm 2, \dots$, $m_{||}$ and $m_{\perp}$ are, respectively, the electron masses corresponding to the longitudinal and transverse motions in the film, Δ_s is the characteristic thickness of the shell, $q_s = 2\pi\nu / \Delta_s$, $\nu = 0, \pm 1, \pm 2, \dots$, f is the Fermi function, and $\delta \rightarrow +0$.

The formula for the longitudinal permittivity of a thin plasma layer Eq. 12.310 differs from that in an unbounded plasma by that the integration with respect to the transverse momentum of an electron is replaced by the summation over the film quantum number n .

Below, we restrict ourselves by the consideration of long-wave plasmons with energies which are small as compared to the energy differences between the energy subzones. That is, we assume that the inequalities

$$q\Delta_s \ll 1, \quad \hbar\omega \ll \varepsilon_{n+1} - \varepsilon_n \quad (12.311)$$

are valid. We consider also that the electron gas in a plasma shell is strongly degenerate. For $\varepsilon_{n, k} \leq \varepsilon_F$ (ε_F is the Fermi level of the current carriers in the film), we replace the distribution function by 1. By integrating in Eq. 12.310

with respect to the component of the longitudinal momentum of an electron which is normal to the vector $\mathbf{q}$, we get

$$\varepsilon^l(q_s = 0) = 1 + \frac{4m_{||}e^2}{\pi\hbar^2\Delta_s q^3} \sum_n k_n \int_{-1}^1 dx \left(\frac{\sqrt{1-x^2}}{x-w_-} - \frac{\sqrt{1-x^2}}{x-w_+} \right), \quad (12.312)$$

where $w_{\mp} = s/\nu_n \mp q/2k_n$, $s := \omega/q$, is the phase velocity of a plasma wave, ν_n is the Fermi velocity, and K_n is the Fermi wave vector of an electron in the n -subzone. After the integration, the longitudinal permittivity takes the form

$$\begin{aligned} \varepsilon^l(q_s)|_{q_s=0} = 1 + \frac{4}{a\Delta_s q^2} \\ \times \left\{ N - \sum_{n=1}^N \frac{k_n}{q} \left[\sqrt{\left(\frac{s}{\nu_n} + \frac{q}{2k_n} \right)^2 - 1} - \sqrt{\left(\frac{s}{\nu_n} - \frac{q}{2k_n} \right)^2 - 1} \right] \right\}. \end{aligned} \quad (12.313)$$

Here, N is the number of filled subzones in the film, and $a = \hbar^2/m_{||}e^2$ is the “effective” Bohr radius.

The analysis of the dispersion equation with permittivity Eq. 12.313 shows that a new type of plasma waves with the acoustic law of dispersion can exist in sufficiently thin films besides the known surface waves with the frequency $\omega \approx \omega_p \sqrt{q\Delta_s/2\varepsilon_0}$, where ε_0 is the permittivity of the medium surrounding a plasma layer. These waves are related to the quantization of the transverse motion of an electron. In thick shells where the film quantization is insignificant, they disappear.

Upon the derivation of the dispersion relation for acoustic waves, we consider that, together with conditions Eq. 12.312, the inequality

$$\left| \frac{s^2}{\nu_n^2} - 1 \right| \gg \frac{q}{k_n} \left(\frac{s}{\nu_n} + \frac{q}{4k_n} \right) \quad (12.314)$$

is valid. In this case, the dispersion equation can be significantly simplified:

$$N - \sum_n \left[1 - \left(\frac{\nu_n}{s} \right)^2 \right]^{-1/2} = 0. \quad (12.315)$$

The influence of the dielectric medium surrounding a layer on the spectrum of acoustic waves turns out to be slight.

Let the phase velocity of a wave lie between the Fermi velocities of the $(m-1)$ -th and $(m+1)$ -th layers, i.e.,

$$\nu_N < \nu_{N-1} < \cdots < \nu_{m+1} < s < \nu_{m-1} < \cdots < \nu_1.$$

We separate s into the real and imaginary parts $s + s_1 - is_2$ and assume that $s_2 \ll s_1$. The phase velocity s_1 is close to the Fermi velocity ν_m . Therefore, we can set $s = \nu_m$ in all terms besides the m -th one. Then Eq. 12.315 takes the form

$$(a + ib) \left[1 - \frac{v_m^2}{s_1^2} - 2i \frac{s_2}{s_1} \right]^{1/2} = 1, \quad (12.316)$$

where $a = N - \sum_{n=m+1}^N \frac{v_m}{\sqrt{v_m^2 - v_n^2}}$, $b = \sum_{n=1}^{m-1} \frac{v_m}{\sqrt{v_n^2 - v_m^2}}$. Then we get, from Eq. 12.316:

$$s_1 = v_m \left[1 - \frac{a^2 - b^2}{(a^2 + b^2)^2} \right]^{-1/2}, \quad s_2 = s_1 \frac{ab}{(a^2 + b^2)^2}. \quad (12.317)$$

As seen from this formula, the decay of the acoustic wave is small (for $\max(a^2, b^2)/(a^2 + b^2)^2 \ll 1$). The most weakly decaying wave has $m = N$. For it, we have $a = N > b$ and $s_1 \sim v_N(1 + 1/2N^2)$, $s_2 \leq s_1/N^2$. With decrease in m , the decay of a wave increases quite rapidly. As for the phase velocity of a wave, it approaches the relevant Fermi one (v_m) and then becomes somewhat less than the latter.

Thus, in thin concentric layers of the electron-nucleus plasma, the ion and electron plasma oscillations have acoustic character, i.e., the frequency depends linearly on the wave vector in the region of small wave numbers:

$$\omega = s_n k. \quad (12.318)$$

In this case, the permittivity can be written in a model form consistent with the model of “jelly” for solids (see Refs. 365 and 372):

$$\varepsilon(\omega, \mathbf{k}) = 1 + \frac{\kappa_d^2}{k^2} + \frac{\kappa_s^2}{k_s^2} - \frac{\omega_{pe}^2}{\omega^2}. \quad (12.319)$$

Here ω_{pe} is the plasma electron frequency of oscillations, $1/\kappa_d$ is the characteristic radius of the screening of nuclei by electrons of a plasma thin layer, $\kappa_s = \frac{2\pi}{\Delta_s}$, k_s is the wave number of a wave along the shell.

In this model, the Fourier component of the interaction energy of nuclei takes the form

$$\begin{aligned} U(\omega, k) &= \frac{4\pi Z_1 Z_2 e^2}{\varepsilon(\omega, \mathbf{k}) k^2} = \frac{4\pi Z_1 Z_2 e^2}{k^2 + \kappa_d^2} \left(1 + \frac{\omega_k^2}{\omega^2 - \omega_k^2} \right) \\ &= U_s(\omega, k) + U_d(\omega, k), \end{aligned} \quad (12.320)$$

where ω_k are the frequencies of longitudinal waves propagating along the system, whose permittivity is given by formula Eq. 12.319. In Eq. 12.320, the first term $U_s(\omega, k) = \frac{4\pi Z_1 Z_2 e^2}{k^2 + \kappa_d^2}$ corresponds to the statically screened Coulomb interaction of two nuclei, and the second term

$$U_d(\omega, k) = \frac{4\pi Z_1 Z_2 e^2}{k^2 + \kappa_d^2} \frac{\omega_k^2}{\omega^2 - \omega_k^2} \quad (12.321)$$

describes the dynamical screening, i.e., the influence of natural oscillations of the medium on the interaction of charges according to the diagram in Fig. 12.67. In Fig. 12.68, we present the Fourier component of the energy of the Coulomb interaction of ions upon their dynamical screening versus the radial wave number.

As seen from Fig. 12.68, the Fourier component of the Coulomb energy has a singular point at a definite relation between the components of the wave vector.

In the interval $0 < \omega < \omega_k$, the interaction changes its sign and becomes the attractive one. In all other regions of the wave vectors and frequencies, the Coulomb barrier turns out to be significantly suppressed like the physical situation arising upon the formation of strongly correlated pairs of electrons in the superconducting state in metals (see Ref. 365) or superfluidity in the nuclear substance upon the interaction of nucleons (see Ref. 292).

In our case, the region of attraction is defined by the eigenfrequencies of a thin plasma layer.

As known, the existence of the acoustic branch of oscillations in a medium enhances the efficiency of the “Cooper” pairing (see Ref. 365). By

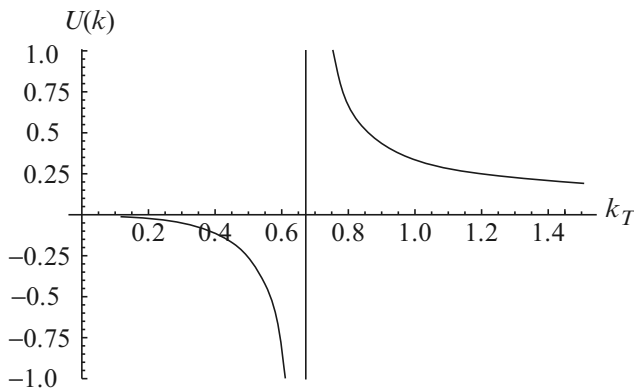


Fig. 12.68. Fourier transform of the interaction potential versus the wave number along the radius.

using the theory of dielectrics, we may estimate the temperature of the phase transition, below which there arises the pairing of identically charged fermions.

In the weak-coupling approximation, the general formula for the temperature of the phase transition as a simple compact functional of the permittivity reads (see Ref. 365)

$$T_c \approx \prod_k (\Omega_k)^{A_k / \sum_k A_k} \exp \left(- \frac{1}{\sum_k A_k \mu} \right), \quad (12.322)$$

where

$$\Omega_k = (\omega_j, E_F), \quad A_k = (a_j (1 + \mu \ln (E_F / \bar{\xi})), -1), \\ \bar{\xi} = \prod_j (\omega_j)^{a_j / \sum_j a_j}, \quad a_j = \left(\omega^2 \partial \varepsilon / \partial \omega^2 \right)_{\omega^2 = \omega_j^2(k)}^{-1}, \quad \omega_j(k)$$

is the discrete collection of the branches of longitudinal oscillations. In the approximation of one mode of oscillations with regard to the dispersion of a thin plasma layer in the “jelly” model Eq. 12.313, the force constant a reads

$$a = \frac{\omega^2(k)}{\omega_p^2} = \frac{1}{1 + \omega_p^2 / k^2 c_{str}^2}. \quad (12.323)$$

For the subsequent estimates, we set $k \approx \bar{k} \approx k_F$. By substituting relation Eq. 12.323 in Eq. 12.316, we get the estimate for the critical temperature: $T_c \approx \hbar \omega_p \approx 8$ keV.

Probability of the Reaction of Synthesis of Nuclei. In the quasi-classical approximation, the tunneling probability across a Coulomb barrier can be presented as a simple integral

$$G_f(\varepsilon, E, n) = \exp \left(- \frac{2}{\hbar} \int_{r_1}^{r_2} \sqrt{2\mu(U - \varepsilon)} dr \right). \quad (12.324)$$

The probability of the reaction of synthesis can be got by the averaging of the tunneling probability with the energy distribution function of nuclei.

Though the nucleus–nucleus collisions and interactions are the very complicated collective process, its description is frequently carried out in the simple approximation of thermal equilibrium. The equilibrium thermodynamics is used, in particular, for the description of the distribution of the products of reactions and the fragmentation of nuclei by momenta. However,

the nonequilibrium thermodynamical states of the nuclear substance are essentially affected in reality by the correlations of the system of nucleons and a formed spatial-field structure in the phase space.

Since the Coulomb barrier in layers of a plasma-field structure, which is formed in the target due to a complicated dynamical process, is strongly suppressed, nucleons can form a cluster in the form of a spherical layer with small thickness almost without obstacles. We note that the conditions of the formation of a nuclear cluster are rapidly improved while moving toward the target center.

The probability of the reactions of synthesis in the regions with extreme parameters is defined by the tunneling probability through a Coulomb barrier G_f with regard to the influence of natural oscillations of the medium on the interaction upon the collision of nuclei with energy ε . We write the total probability of the reaction of synthesis, by integrating the tunneling probability through the barrier over all energies, as

$$P_{tot}(E, n) = \int_0^{\infty} G_f(\varepsilon, E, n) \sqrt{\frac{2\varepsilon}{\mu}} f_q(\varepsilon) d\varepsilon. \quad (12.325)$$

In Fig. 12.69, we present the probability of the reaction of synthesis Eq. 12.325 in a nonequilibrium state as a function of the parameter of nonextensiveness q .

Thus, the compensation of the Coulomb field of nuclei, nonequilibrium states of the system, and collective polarization effects, being a consequence of the dynamical processes running in a target under the action of a pulse action, give a possibility to join the nuclei from the layers with enhanced density into a single nuclear system.

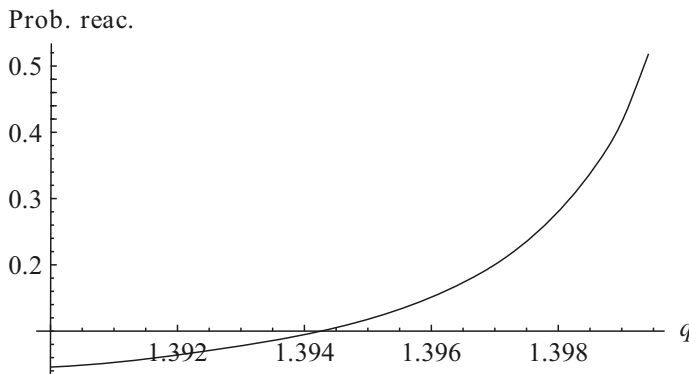


Fig. 12.69. Probability of the reaction of synthesis of nuclei versus the parameter of nonextensiveness q .

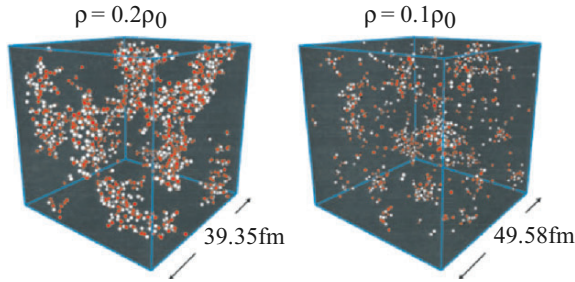


Fig. 12.70. Nuclear clusters in the form of Nuclear Pasta (see Ref. 367).

We recall that these layers belong to the plasma-field structure which arises only under conditions of the development of self-consistent two-fluid dynamical processes in the electron-nucleus plasma of a target with regard to the interaction of the flows of electrons and nuclei.

The appearance of the collective attraction between nuclei in the thin concentric spherical layers of a target (i.e., in the system of shells) induces the process of growth of a cluster consisting of the nucleons of nuclei of the target.

Last years, the growth of clusters in the nuclear substance has received a lot of attention. A number of works (Refs. 373, 374) was devoted to the numerical simulation of the processes of growth of fractal clusters in the nuclear substance with the use of the Monte-Carlo quantum method. As a result, the appearance of complex structures, which were named Nuclear Pasta, was established. The origin and sense of this name are clearly seen from Fig. 12.69, where the results of the simulation of 2048 interacting particles in a cell are presented.

We could imagine the dynamics of the growth of clusters by solving the chain of equations for the distribution function of clusters of various sizes (see, e.g., (see Ref. 277)) or the system of equations which are similar to the equations of chemical kinetics and follow from the kinetics of clusters after some simplification. However, this program is too labor-consuming. Therefore, we will estimate the dynamics of the formation of a global cluster (a cluster containing a macroscopic number of nuclei) from elementary clusters (nuclei) by essentially simplifying the system of equations with preservation of their form. In this case, the system of equations is reduced, in fact, to the Smoluchowski equation.

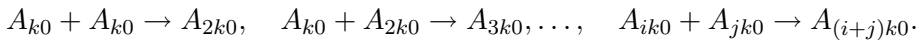
Formation of Nuclear Clusters. The growth of a cluster composed from elementary clusters can be described by the kinetic equation (see Ref. 355) proposed by Smoluchowski for the description of coagulation in

1914. From the very beginning, this equation is used for the description of the phenomena of aggregation in various fields of science.

Consider a system of aggregating particles, in which the formation of clusters of various sizes continuously occurs as a result of the binary contacts of particles, as well as formed clusters, with their subsequent union.

For the statistical description of macroscopic properties of the system, it is necessary to know the distribution of clusters by size. As the size of a cluster, we take the number of nucleons in it. So, we will seek for the time behavior of the concentrations C_k of clusters consisting of k nucleons.

Thus, we want to describe the consistent system of reactions



An equation for the concentration C_k can be written in the form of a Smoluchowski coagulation equation

$$\frac{dC_k}{dt} = \frac{1}{2} \sum_{i+j=k} K_{ij} C_i C_j - C_k \sum_{j=A}^{\infty} K_{kj} C_j. \quad (12.326)$$

As the initial condition for this equation, we take the state, in which all nuclei are of the same size k which is set for the sake of clearness as the mass number A equal to that of the target substance, i.e., $C_k(0) = \delta_{kA}$. For the probability K_{ij} of the fusion of clusters with the sizes i and j , we may use the approximation, in which this probability is proportional to the product of the surface areas of the initial clusters: $K_{ij} \propto (ij)^{2/3}$.

In such an approximation, the Smoluchowski equation can be integrated analytically (see Ref. 355), and the analysis of its solution shows that the mean size of a cluster can become infinite for a finite time, namely the duration of the phase transition into the state of gel.

The dependence of the mean size of a cluster on time by this solution is presented in Fig. 12.71 (a, b), where we give the time in units of the phase transition duration.

As seen, the plot of this dependence has form typical of the explosive instability. For a finite time, there occurs the transition of a nuclear cluster in the gel phase, and all monomers tend to form one cluster (in the limit, it is infinitely large). The rate of this phase transition is high, and, therefore, it is convenient to consider the process of growth on the logarithmic scale (see Fig. 12.71 [b]). Due to the phase transition, the nuclear cluster of a macroscopic size appears.

We note that, as is known (Ref. 290), fractal clusters can be formed upon the clusterization. Their peculiarity is the very great surface, being much greater than the surface of a spherical layer, in which they are being

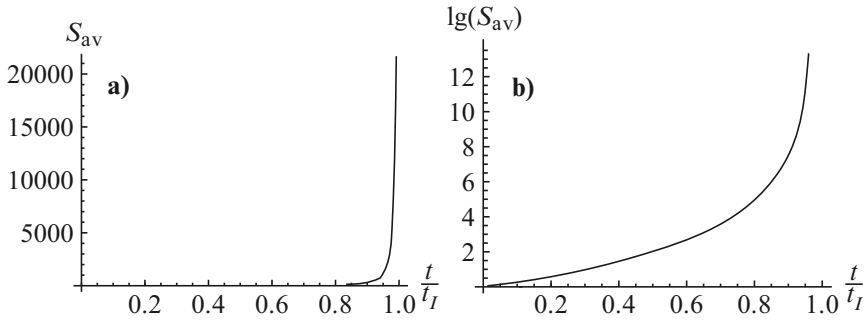


Fig. 12.71. Mean size of a cluster (a) and its logarithm (b) versus the dimensionless time.

formed. This surface, hence, has great surface energy, which allows a macroscopic system of fractal clusters to play the role of an extraordinarily efficient accumulator of a huge amount of energy.

Thus, under conditions of the development of the nonlinear stage of the evolution of a spatial-charge wave, a great number of thin layers appears. These layers become, as a result of the phase transition, fractal nuclear clusters. The fractality of a cluster means that it possesses a complicated spatial structure.

As indicated above, the possibility for fractal structures to appear was earlier considered in many works (see, e.g., Refs. 219, 280) on the basis of the analysis of phase transitions in the nuclear substance and the experimental data on the collisions of heavy ions.

The densities of a substance and a charge are distributed in a cluster by the law with power asymptotics (see Ref. 290):

$$\rho_Z(r) \approx \frac{\text{const}}{r^{d-D_f}} = \frac{\text{const}}{r^\alpha}, \quad (12.327)$$

where d and D_f are, respectively, the topological and fractal dimensionalities of a cluster (see also Sec. 12.2).

By the estimates of the growth of 3D clusters ($d = 3$) under a great probability of adhesion, the fractal dimensionality of a cluster $D_f = 2.39$, and $\alpha = 0.61$ in Eq. 12.327. We will describe the density distribution of the nuclear substance in a cluster by starting from the idea of the existence of strong correlations between values of the density at different points of the nucleus, which depend by a power law on the distance between them, under the formation of nuclear structures.

As usual, the density distribution of a substance in a continuous nucleus in the form of a spherical drop is approximated with the use of the Fermi function $\rho_{0A}(r, R, \delta) = \frac{1}{1 + \exp(\frac{r-R}{\delta})}$ (see Ref. 375). Here, R is

the characteristic radius of the nucleus, and δ is the characteristic size of a diffusion zone. For a fractal cluster with regard to correlations, the density distribution can be approximated by the following function defined by the Tsallis quantum distribution Eq. 12.298:

$$\rho_A(r, R, \delta, q) = \frac{\theta(R - r)}{1 + \exp_q\left(\frac{r - R}{\delta}\right)} + \frac{\exp_q\left(-\frac{r - R}{\delta}\right)}{\exp_q\left(-\frac{r - R}{\delta}\right) + 1} \theta(r - R). \quad (12.328)$$

Here, the parameter q is expressed via the fractal dimensionality D_f of a cluster: $q = 4 - D_f$. The density distribution functions satisfy the relation $\rho_A(r, R, \delta, q \rightarrow 1) = \rho_{0A}(r, R, \delta)$.

In Fig. 12.72, we display the distribution functions for both a continuous nucleus and a fractal cluster of the same size for $A = 200$ with the parameters corresponding to a random process of adhesion of new nuclei to the cluster ($D_f \approx 2.39$).

As shown in Ref. 376, with the use of the theory of Hartree–Fock–Bogolyubov with the effective interaction D1S Gogny (see Ref. 377), the inhomogeneities of the distribution of the nuclear substance in a nucleus (see Fig. 12.72) lead to the possibility for superheavy nuclei to exist in a wide range of mass numbers. In Ref. 376, two types of “superheavy nuclei” were discovered: true “bubbles” with actually vanishing nuclear density in the central region of the nucleus and “quasibubbles” (“unsaturable nuclei”) with a lesser but finite density near the nucleus center (see Figs. 12.73 and 12.74).

For quasibubbles and true bubbles, the nuclei were stable for mass numbers A (and numbers of protons Z) in the ranges, respectively, $292 <$

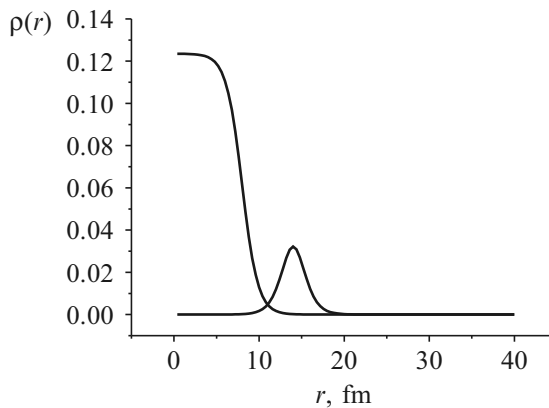


Fig. 12.72. Density distribution in a nucleus in the form of a liquid drop and in a nucleus with power spatial correlations of the density in the form of a bubble corresponding to the idea of a nucleus as the object with a complicated spatial structure.

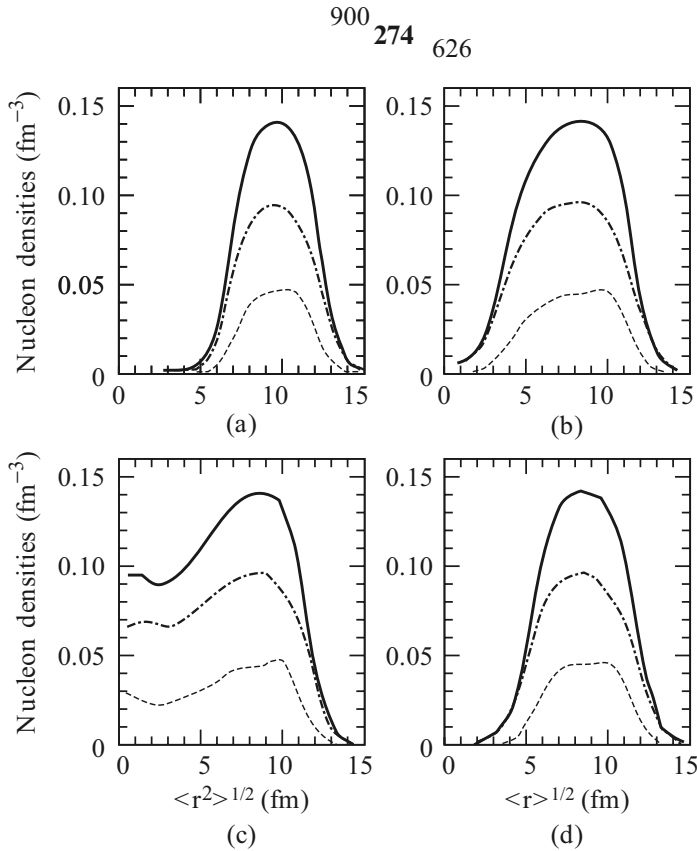


Fig. 12.73. Distributions of the density for bubble and quasibubble nuclei. The distribution of protons (*dashed curves*), neutrons (*dash-dotted curves*) and the total nucleon density (*solid curves*).

$A < 750$ ($120 < Z < 240$) and $750 < A < 920$ ($240 < Z < 280$) (see Fig. 12.74). The dominant way of the decay of quasibubbles is the α -decay, whereas the true bubbles decay through the reaction of fission.

The characteristic lifetimes of quasibubble nuclei relative to the α -decay have values from several seconds to microseconds, and the characteristic times for the reaction of fission are in the range from several years to microseconds.

The spatial structure of such superheavy nuclei can be studied better, if we pass to the hydrodynamic level of the description of the nuclear substance and consider the stationary flows of various components of a nucleus (in particular, electrons and nucleons). Electrons involved in a flow are formed as a result of the condensation of electron-positron pairs from vacuum in the field of nucleons.

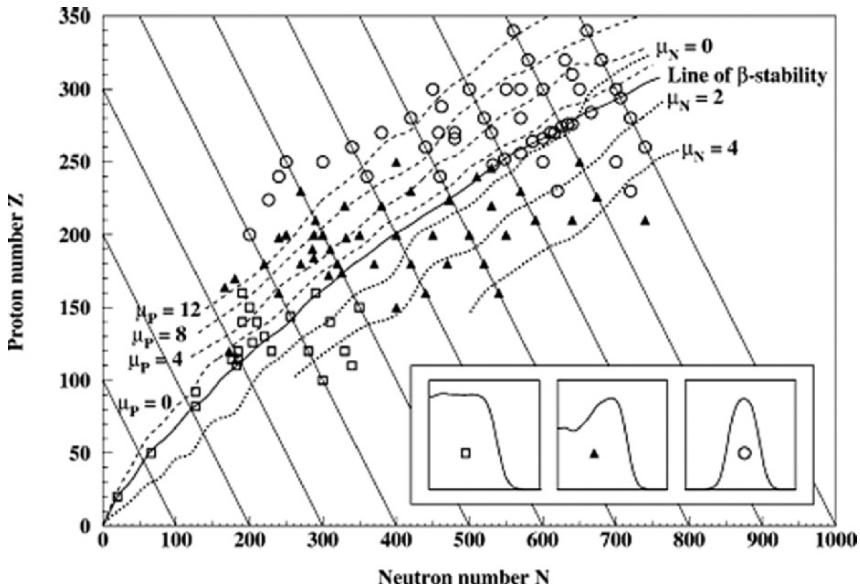


Fig. 12.74. Stability lines of bubble and quasibubble superheavy nuclei.

Besides electrons, other particles can be created from vacuum in the strong electrostatic fields inside nuclei. The properties of nuclei are most significantly affected by the processes of condensation of negatively charged bosons (π^-) from the corresponding pairs of mesons created in the field of a nucleus from vacuum and by the inhomogeneity of their distribution in a nucleus.

We note that the revealed nuclear structures influence significantly both the properties of a substance on the macroscopic level upon the formation of plasma-field structures from the system of concentric shells under study and those in a nucleus in the form of a strong inhomogeneity of the distribution of nucleons in a nucleus.

The created system of macroscopic shell structures in the electron-nucleus plasma is not immovable. *The motion of the system of concentric nuclear structures toward the target center is the propagation of a wave involving the self-organizing nuclear combustion and the synthesis of nuclei with energy release.*

The wave with self-organizing nuclear combustion is basically different from a wave with combustion in the inertial synthesis.

Some differences are seen at once, even after a brief analysis. In the inertial synthesis, the wave with nuclear combustion arises on the final stage after the ignition of thermonuclear reactions at the target center as a result of the creation of the sufficiently great density and temperature for the

beginning of the reactions of synthesis in direct binary collisions of the initial nuclei. As seen, the wave with nuclear combustion arises in the self-organizing synthesis on the way toward the center and is accompanied by the growth of the density and the acceleration of the motion of the structure.

The main distinctive feature of the self-organizing synthesis is its collective character. As follows from the above-presented, the reaction of collective synthesis can be schematically presented as

$$\sum_{i,j}^{N_{tg}} \left\{ \begin{array}{l} \left(\begin{smallmatrix} A_{tg} \\ Z_{tg} \end{smallmatrix} N \right)_i + q \rightarrow \left(\begin{smallmatrix} A_{tg} \\ Z_{tg} \end{smallmatrix} N \right)_i^* + q^* \\ \left(\begin{smallmatrix} A_{tg} \\ Z_{tg} \end{smallmatrix} N \right)_j + q^* \rightarrow \left(\begin{smallmatrix} A_{tg} \\ Z_{tg} \end{smallmatrix} N \right)_j^* + q \end{array} \right\} \Rightarrow M^*, \quad (12.329)$$

$$M^* \Rightarrow N_k Fe + \sum_i A_i + N_p p + N_n n.$$

As a result of the establishment of a great number of bonds between the pairs of nuclei of the target $\begin{smallmatrix} A_{tg} \\ Z_{tg} \end{smallmatrix} N$ with the help of the interaction via the exchange by plasmons, there appears a macroscopic nuclear structure M^* (the upper row). Upon the motion toward the center (which is accompanied by the growth of the density of M^*), the formation of new nuclei from the external surface of the macroscopic cluster occurs.

Such a scheme of the synthesis of nuclei is, from our viewpoint, the most promising generalization of the existing models of the nucleosynthesis of heavy elements in stars which are briefly surveyed above.

Prior to this model of self-organizing nucleosynthesis, the most successful theory of the nucleosynthesis of heavy elements (it cannot explain, however, many peculiarities of the abundance of elements) was the theory of the formation of elements as a result of the capture of neutrons by the initial light nuclei from the external reservoir of neutrons (see Refs. 214, 378).

As distinct from the theory of thermonuclear synthesis, this theory uses the assumption that the formation of elements can occur not obligatorily at huge thermonuclear pressures and temperatures, but with the help of a gradual addition of neutrons. However, the “building material” for new nuclei is too uniform (only neutrons) in this approach. In the case of self-organizing nucleosynthesis, the “building material” is the full collection of nucleons of a macroscopic cluster in the amount which is self-consistent for the formation of new elements. The formation of new nuclei occurs upon the evaporation of the external surface of the macroscopic nuclear cluster.

Evaporation of a Nuclear Cluster. Firstly, the statistical ideas were introduced in nuclear physics by N. Bohr in his theory of intermediate nuclei (see Ref. 381), V. Weisskopf in the model of evaporation (see Ref. 382), and L. Landau in the hydrodynamic theory of multiple creation of particles

in collisions. It is natural that great nuclear systems undergo the process of fragmentation or evaporation. These processes occur usually so that the system has time to pass into an equilibrium state prior to the fragmentation.

The distribution of fragments $f(A)$ by the number of nucleons can be calculated within the theory of the “gas–liquid” phase transitions in the nuclear substance and turns out to be a power one (see Ref. 280):

$$f(A) \propto A^{-\alpha_f}. \quad (12.330)$$

In agreement with the thermodynamical theories, we can get the relations for the index α_f by using the Fisher drop model (see Ref. 284). The system of evaporating nuclei and the surface, from which they are evaporating, can be assumed to consist from nuclear drops of various sizes. The statistical sum of a grand canonical ensemble in the Fisher model reads

$$f(A) = f_0 \exp \left(\frac{\mu_1 - \mu_2}{T} A - \frac{4\pi r_0^2 \sigma_A}{T} A^{2/3} - \tau_f \ln(A) \right), \quad (12.331)$$

where τ_f is one of the Fisher critical indices. The Fisher indices are related to the ordinary critical indices of the phase transition point α_{cr} , β_{cr} , γ_{cr} , ν_{cr} with the help of the relations

$$\tau_f = 2 + \frac{1}{\delta_{cr}} = 1 + \frac{d\nu_{cr}}{\beta_{cr}\delta_{cr}}, \quad \sigma_f = (\beta_{cr} + \gamma_{cr})^{-1} = \frac{1}{\beta_{cr}\delta_{cr}}. \quad (12.332)$$

From the viewpoint of thermodynamical theory, the process of evaporation of a macroscopic cluster can run through the gas–liquid phase transition or the appearance of metastable states.

In the mean-field approximation, the critical indices of the gas–liquid phase transition are $\beta_{cr} = 1/2$, $\gamma_{cr} = 1$, $\delta_{cr} = 3$, and $\tau_f = 7/3$.

For the phase transition with the formation of metastable states of nuclear drops or bubbles, the Fisher model gives the critical index of the phase transition $\delta_{cr} = 2$ and, hence, the Fisher critical indices calculated by formula Eq. 12.332 turn out to be $\tau_f = 2.5$, $\sigma_f = 1$.

The Fisher critical index σ_f defines the dependence of the surface area of a cluster on the number of particles, $S \propto A^{\sigma_f}$. Thus, in the phase transition under consideration, the surface area of the cluster $\propto A$, which testifies to the fractal character of the structure. The above-presented estimates yield that $\gamma = \sigma_f \leq 1$ according to Eq. 12.119, and the fractal dimensionality of a cluster $D_f = \frac{2}{\gamma - k} \geq 2.71$ according to Eq. 12.120.

It follows from Eq. 12.331 that the equality $\alpha_f \approx \tau_f$ holds near the phase transition (where $\mu_1 = \mu_2$, $\sigma_A \approx 0$). In view of the above-presented

estimates of the critical indices, we get that the index α_f of the distribution of clusters by the mass number (see Eq. 12.330) is in the interval $2.33 < \alpha_f < 2.5$.

The main contribution to the distribution by A Eq. 12.330 is given by the most stable nucleus at a given shell density. But, the right-hand side of reaction Eq. 12.329 can involve, besides the most probable nucleus, the whole spectrum of nuclei with charges Z_i in the interval $1 < Z_i \leq Z_{\max}$ with lower probabilities.

Upon the motion of the whole structure of a macroscopic cluster toward the target center, the mass numbers of the most stable nucleus and the entire collection of nuclei evaporating from the cluster surface, which are most stable and possess a minimum energy in the general case, vary with increase in the density of the shell substance. With increase in the density of baryons in the shell, most efficient is the evaporation of more and more heavier nuclei. In Sec. 12.1., we have analyzed the thermodynamical equilibrium of free neutrons, nuclei, and electrons and have got the dependence of the mass number of the most stable nucleus on the density of nucleons which is drawn in Fig. 12.14.

Since the most stable nuclei with the least energy at not very high densities are nuclei of Fe (see Fig. 12.75), it is the evaporation of nuclei of Fe that begins the process of self-organizing nuclear synthesis.

As a result, the beginning of the active phase of the evolution of the shell is characterized by the target substance density (most frequently, we executed the experiments with Cu) before the shell to be greater than the density of Fe formed behind the leading edge of the wave with nuclear combustion.

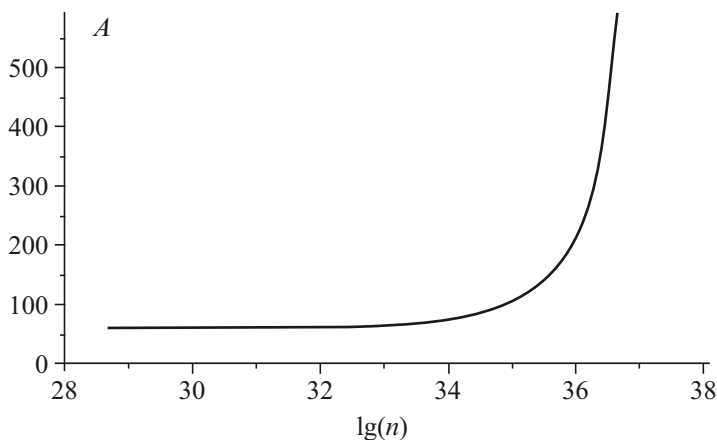


Fig. 12.75. Mass number A of the most stable nucleus versus the density of baryons (on the logarithmic scale).

The difference of the substance densities before the moving shell and behind it leads to that a part of nucleons from the target, which were included in the composition of a cluster, remains in the cluster and participates in the motion toward the target center. The other part is evaporated and remains in the condensed state behind the trailing edge of the moving structure. The energy, which is released upon the evaporation of nuclei, causes the accelerated movement of the evaporating macroscopic shell structure toward the target center.

In the general case, a macroscopic amount of nuclei is evaporated from the cluster surface. Therefore, the reaction scheme including only the initial and final nuclei can contain the number of variants which is by many orders more than that in the case of binary reactions.

Because the collective reaction involves the great number of nuclei of Cu, it is a statistical process.

The statistical character of the process allows the system to maximally use nucleons in their transformation in Fe (or in other nuclei with the maximum stability depending on the shell substance density) with a minimum number of unutilized nucleons and, therefore, to realize the evolution of the system along the most energy-gained trajectory in the phase space.

Since the distribution of evaporated nuclei by mass numbers (charges) and the dynamics of the shell evolution are tightly connected, they are defined by the single variational principle for the whole dynamical process. As such an evolutionary principle, we may consider the self-consistent optimization of the process of nuclear combustion of the initial nuclei A_{tg} through the formation of a megacluster and its evolution according to the conditions of normalization, integral criteria (see Chapter 2), and the flow of nucleons of the target onto the internal surface of the target.

For the further estimates, it is necessary to analyze the binding energy of the nuclear structures arising upon the self-organizing nucleosynthesis.

12.3.3. Binding Energy of Nuclear Structures: A Generalization of the Weizsäcker Formula

During the last decade, the development of nuclear physics is characterized by a rapid growth of the interest in nuclear structures. Structures of the nuclear substance are being studied in ordinary nuclei and in some exotic nuclear objects. The analysis of both the experimental data derived at the Electrodynamics Laboratory “Proton-21” and the above-presented theoretical models allows us to classify the nuclear structures possessing a number of specific properties and to explain the majority of experimentally observed effects.

It has been shown in the previous section that, as a result of the dynamical evolution of nonlinear waves in the electron-nucleus plasma of a target, there appear nuclear clusters with macroscopic sizes. As known, a nucleus can be considered as a system of interacting nucleons, pions, and isobars. At a great density of the system and in the presence of external actions, we cannot neglect the correlations in the system of nucleons involving a complex spatial structure of nuclear clusters. The latter can be described in the framework of various models of the nuclear substance. One of the efficient means to account the correlations present in a nuclear system can include the use of a nonextensive distribution function instead of the Fermi distribution function. Moreover, in the description of the spatial structure of clusters, we can use the notion of fractal dimensionality and the distribution of the mean density as a function of the radius with a power asymptotics, which corresponds to the power decay of the correlations of the density in a cluster.

Macroscopic nuclear systems arisen as a result of the dynamical processes have specific properties which are related, first of all, to the peculiarities of the dependence of their energy on parameters.

Below, we list the main nuclear structures which are of importance in our model:

- Nuclei in the form of a liquid drop in the standard model
- Nuclei as liquid drops with a diffusive boundary and strong spatial correlations
- Bubble nuclei, i.e., nuclei in the form of a bubble with a thin shell
- Bubble nuclei with a fractal shell
- Macroscopic fractal shell
- Macroscopic nuclear structure in the form of a system of concentric shells

Consider the simple model of a nucleus containing A nucleons (Z protons and $N = A - Z$ neutrons) as a nuclear drop of the Fermi liquid. In this case, the total energy of a nucleus includes:

- energy of the Coulomb field of Z protons (W_q)
- energy of the nucleus surface (W_s)
- energy of the nuclear substance in the nucleus volume (W_v) with regard to the symmetry-related energy associated with the processes of β -decays and the balance of neutrons and protons in the nucleus (W_a)
- energy of the condensates of various particles in the nucleus ($W_\pi, W_e, \dots$)
- energy of a self-consistent electromagnetic field in the dynamical nuclear structure (W_{em})

- kinetic energy of the stationary hydrodynamic flows of the nuclear substance in dynamical nuclear structures (W_{gd})

The energy of the self-consistent electromagnetic field and the energy of the hydrodynamic flow of the nuclear substance are included in the binding energy of the nuclear structure, because they ensure the existence and stability of nuclear structures, as well as the strong correlations in the nuclear substance through long-range electromagnetic fields.

As an example of such macroscopic nuclear structures, we mention the macroscopic nuclear fractal structures arising as a result of the clusterization on the leading edge of a nonlinear wave of the volume charge density (see Sec. 12.2 and vortex structures in a nucleus which are similar to the structures arising in the nuclear substance (see Ref. 300).

It is commonly accepted that all components of the binding energy depend on the density of the nuclear substance ρ , mass number A , and charge of a nucleus Z (the number of protons in a nucleus).

From six above-listed components of the nucleus energy, three first ones are terms in the well-known Weizsäcker formula. The term related to the processes of condensation was added to the Weizsäcker formula in Migdal's works (see Ref. 300).

Below, we will estimate the energy of the self-consistent field in nuclear clusters for the case of a fractal structure described in Sec. 12.2. The above-discussed plasma-field structures can be generalized on the basis of the analysis of the hydrodynamic flows and the self-consistent electromagnetic fields in a nucleus. Such a generalization allows us in future to refine the structure of a cluster on its surface. However, we think that this refinement does not lead to any qualitative changes in the binding energy.

It is clear that, in the case of great nuclear clusters, the binding energy will additionally depend on 4 parameters: q , D_f , n_e , and E which are lacking in the traditional Weizsäcker formula. Two first parameters characterize the cluster itself. They are the parameter of nonextensiveness q of the nuclear substance (the degree of correlation of states in a nucleus) and the fractal dimensionality D_f of a nuclear cluster (or the parameters related to it). Two last parameters correspond to the macroscopic properties of the medium, where nuclear clusters were formed. That is, D_f and q should depend, first of all, on the density of electrons and the intensities of components of the collective electromagnetic field.

Let us now pass to the analysis of separate terms in the energy of a nuclear cluster.

Processes of Condensation in a Nucleus. The idea proposed and developed by A. Migdal consists in that the loss of the energy of the baryon

subsystem upon the growth of the density can be compensated by a gain at the expense of the appearance of the π -condensate. In this case, the process of compression of the system will proceed until there occurs the phase transition into a new quasistationary state with positive compressibility.

As known, Bose-particles are described by the Klein–Gordon equation. For the states without external field, it has the form

$$\Delta\psi + (\omega^2 - 1)\psi = 0. \quad (12.333)$$

If a state has a definite momentum k , then $\omega = \pm\sqrt{1 + k^2}$. Two signs correspond to a particle and an antiparticle, and Eq. 12.333 describes them simultaneously.

In a nucleus, the positive charge of protons creates the field, where electrons and negative pions can be generated. Positrons, neutrinos, and mesons, which are created together with electrons and negatively charged pions, rapidly leave the system. These processes are described by the Lagrange function density for the pionic and electric fields:

$$L = [(\omega - V)^2 - 1]|\varphi|^2 - \frac{\Lambda|\varphi|^2}{2} + \frac{(\nabla V)^2}{8\pi e^2} + \rho_p V - |\nabla\varphi|^2, \quad (12.334)$$

where ρ_p is the proton charge density. The variation in φ and V gives the system of equations defining the fields:

$$\Delta\varphi + [(\omega - V)^2 - 1]|\varphi|^2\varphi = 0, \quad \Delta V = 4\pi e^2 [\rho_p - 2(\omega - V)|\varphi|^2]. \quad (12.335)$$

The simple estimates for this system can be performed under the assumption of the quasihomogeneity of the positive charge distribution over the nucleus. Then the densities of electrons and negative mesons are

$$\rho_{e,\mu} = -(1/3\pi^2) [(\varepsilon_{\max} - v)^2 - m_{e,\mu}^2] \theta(|V(r) - \varepsilon_{\max}| - m_{e,\mu}). \quad (12.336)$$

Here, $\varepsilon_{\max} = \mu_n - \mu_p$ is the limit energy, $\mu_n = (\partial E(N, Z)/\partial N)|_Z$ and $\mu_p = (\partial E/\partial Z)|_N$ are, respectively, the chemical potentials of neutrons and protons, $E(N, Z)$ is the total energy of the system, and $|\varepsilon_{\max}| < m_e$ (due to β -processes). The contribution to the energy per particle, which is defined by the π -condensate charge, has the form $\varepsilon(\rho_\pi) = 2\rho_\pi^2/F_\pi^2\rho$, and the electric charge density of the π -condensate is defined by the condition of local electroneutrality,

$$\rho\partial\varepsilon(\rho_\pi)/\partial\rho_\pi - V = \text{const.} = 0. \quad (12.337)$$

From here, we get the electric charge density for the hadronic subsystem as

$$\rho_h = \left(\frac{\rho}{2} + \frac{F_\pi^2 V}{4}\right) \theta(R - r). \quad (12.338)$$

The electric potential satisfies the Poisson equation (see Eq. 12.335). At $Ae^3 \ll 1$, electrons play a minor role, but the inhomogeneity of the pionic charge distribution is significant. The solution of the Poisson equation reads

$$V = - \begin{cases} \frac{2\rho}{F_\pi^2} + \frac{C}{r} \operatorname{sh} \frac{r}{l}, & r < R, \\ \frac{Z_\infty e^2}{r}, & r > R, \end{cases} \quad (12.339)$$

where $l = (F_\pi e \sqrt{\pi})^{-1} \approx 5m_\pi^{-1}$, $C = -2\rho l / F_\pi^2 ch(R/l)$. From here, we can find the charge observed at infinity as

$$Z_\infty \approx \frac{3}{2} \left(\frac{l}{R} \right)^3 \left(\frac{R}{l} - \operatorname{th} \frac{R}{l} \right) A. \quad (12.340)$$

It is seen from Eq. 12.343 that the effective charge of a nucleus depends essentially on its mass number. Such nontrivial dependence of the effective charge on A leads to a change of the β -stability line as compared to the traditionally used formula $Z_{\beta opt} \approx \frac{A}{1.9+0.015 A^{2/3}}$.

For $R \gg l$, the pionic charge distribution in a nucleus is strongly rearranged even without regard for the electron screening. The nucleus interior becomes electroneutral. The charge observed at infinity has a simple asymptotics:

$$Z_\infty \approx 2^{-1/3} (\rho/\rho_0)^{2/3} A^{1/3} / F_\pi^2 e^2 \ll A. \quad (12.341)$$

In the limit case where $Ae^3 \gg 1$, we can find the analytic solution of Eq. 12.332, by assuming

$$V \approx -V_0 \chi(x), \quad \chi(x) = \begin{cases} \beta/(x+b), & r > R, \\ 1 - C \exp(\mu x), & r < R, \end{cases} \quad (12.342)$$

where $x = (r - R)/l$, $l = 2eV_0/\sqrt{3\pi}$; b , C , μ , β are constants; $\chi(-\infty) = 1$, and $\chi(\infty) = 0$.

The value of V_0 can be found from the condition of full electroneutrality inside the system. Let us introduce the density $\rho^* = \sqrt{3/2} (\pi/2) F_\pi^2 \approx 10\rho_0$, at which a half positive charge of a nucleus is compensated by fermions (e^- , μ^-), and a half by π -mesons of the condensate. Then, in the limiting case $\rho \ll \rho^*$, we get

$$\beta \approx \sqrt{2}, \quad V_0 \approx \frac{2\rho}{F_\pi^2}, \quad \mu \approx \sqrt{3 + 3\pi^2 F_\pi^2 / 4V_0^2}. \quad (12.343)$$

At $\rho \ll \rho^*$, a significant contribution to the charge screening inside a nucleus is given by π -mesons of the condensate. Outside a nucleus, the screening is realized by electrons.

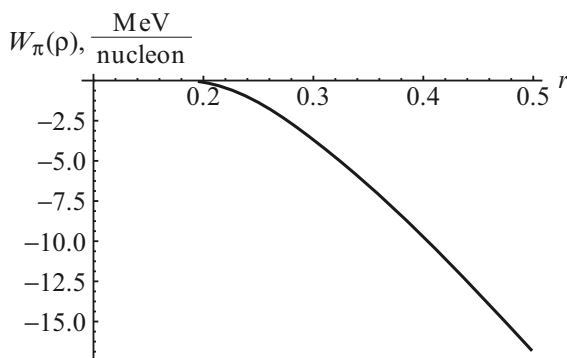


Fig. 12.76. Condensation energy versus the nuclear substance density.

Thus, in the limit case of a great superheavy nucleus [the corrected criterion looks as $Ae^3 \gg (\rho^*/\rho)^2$], we deal with the electroneutral plasma composed of baryonic quasiparticles, condensed pions, electrons, and negative muons.

In estimates, the approximation of a homogeneous nuclear substance is usually used. However, the account of both the accumulation of fermions and the nonuniformity of the distribution of the pionic charge improves the conditions for the stability of anomalous nuclei at the same values of nuclear constants as without regard for these effects. Thus, to correctly describe the stable nuclei, it is necessary to consider the inhomogeneity of the distribution of all components of a nucleus.

The contribution of the condensate to the energy of a nucleus is shown in Fig. 12.76. This plot corresponds to a dependence on density approximated by the relation

$$W_\pi(\rho) = -70 \frac{\wp_E(E_f, \beta, q)}{\wp_E(E_f, \beta, 1)} \left(A + B\rho/\rho_c + C(\rho/\rho_c)^2 \right) (\rho - \rho_c)^2/\rho + \alpha_\pi (1 - 2Z/A)^2, \quad (12.344)$$

where (see Ref. 293) $\rho_c = 0.185$, $A = 1.26$, $B = -0.08$, $C = -0.07$, and $\alpha_\pi = 19.6\rho/\rho_0$.

Contribution of the Surface to the Binding Energy of a Nucleus.

As usual, the contribution of a surface to the binding energy is estimated under the assumption of the quasiumiformity of the nuclear substance in a cluster. Let us consider explicitly the inhomogeneity of the substance distribution in a nucleus upon the estimate of its surface energy.

In the quasiclassical approximation, the dependence of the surface energy on the density can be presented as

$$W_s(A, D_f, \rho, \delta) = 4\pi \int \lambda_N (\rho_A) (\nabla \rho_A)^2 r^2 dr. \quad (12.345)$$

In this formula, λ_N is a phenomenological parameter defined by the effective radius of the nucleon–nucleon interaction, and $\rho_A(r)$ is the density of the nuclear substance. This gives also the ordinary formula for the surface energy of a nucleus considered as a liquid drop with radius R and with the thickness of the fuzzy edge of a nucleus δ :

$$W_s \approx 4\pi\lambda_N \left(\frac{\rho}{\delta}\right)^2 \frac{1}{3} (R^3 - (R - \delta)^3) \approx \left(\frac{1}{3}\lambda_N \frac{\rho^2}{\delta}\right) 4\pi R^2 \approx \sigma_s 4\pi R^2. \quad (12.346)$$

Taking formula Eq. 12.328 for the nuclear substance density, we can deduce the dependence of the surface energy on the radius and the mass number from relation Eq. 12.344, including the nuclei with strong correlations of their density. In this case, the absolute value of the surface energy decreases. For $q = 1.61$, its dependence on the mass number is shown in Fig. 12.77 together with the traditional dependence for a liquid drop.

For such correlations of the density, the stability of nuclei enhances.

As usual, nuclei and nuclear clusters are assumed to have form close to that of a sphere. By virtue of the discussed mechanisms of formation of nuclear clusters, further we will consider also nuclei in the form of a spherical shell with the external radius R , internal radius R_1 , and thickness δ ($\delta = R - R_1$).

If a nuclear cluster appears in the form of a spherical layer with a homogeneous distribution of the nuclear substance in it, the following

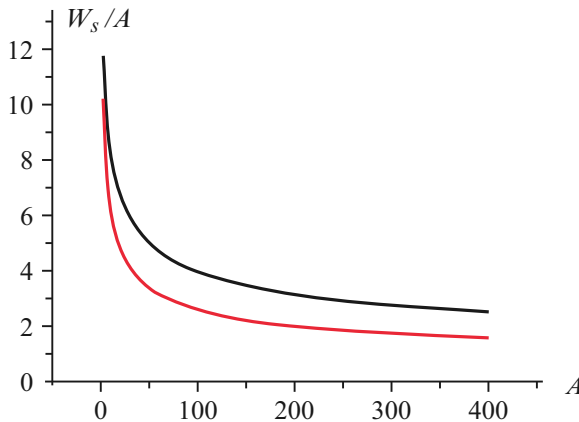


Fig. 12.77. Surface energy versus the mass number for nuclei in the form of a liquid drop (upper curve) and nuclei with power spatial correlations.

relation between the external radius cluster R , layer thickness δ , and mean density ρ_0 is valid: $R^3 - (R - \delta)^3 = 3\delta R^2 - 3\delta^2 R + \delta^3 = (3/4\pi\rho_0)A$ or

$$R_\delta(A, \rho, \delta) = \sqrt{\frac{A}{4\pi\delta\rho} - \frac{\delta^2}{12}} + \frac{\delta}{2}. \quad (12.347)$$

The plot of this dependence is given in Fig. 12.78(a, b), respectively, for the ordinary density of the nuclear substance and for the density by two orders less.

The analysis of these figures yields that nuclear clusters will have macroscopic sizes if their mean density is much less than the ordinary density of nuclei.

This corresponds to the conclusions about the porosity of a cluster and about the possibility to form a cluster with a complicated structure which were presented above upon the consideration of the clusterization in plasma-field structures.

The fractal structure of the surface of nuclear systems with great size causes, naturally, an increase in the surface energy $W_s(A, D_f, \rho, \delta)$ and a change of its dependence on the mass number, which should be accounted in the formula for the binding energy instead of the traditional term $W_s = c_2 A^{2/3}$.

Below, we give some estimates for the surface energy of a fractal cluster.

In view of the relations derived in Sec. 12.1., we get:

$$W_s(A, D_f) = c_2 A^{2/3} (A/A_0)^{\gamma-2/3}, \quad c_2 = 4\pi\sigma_0 = 17.8, \quad (12.348)$$

where A_0 is the mass number of the elements, from which the cluster is built.

The dependence of the surface energy per nucleon on the mass number for the limit values of the fractal dimensionality is presented in Fig. 12.79.

The lower curve corresponds to a smooth surface, and the upper one does to a maximally porous surface.

Energy of the Strong Interaction of the Nuclear Substance. The energy of the strong interaction of the nuclear substance can be estimated by the formula

$$W_v(\rho, \beta, q) = 140 \left(-0.1125 + \frac{0.14 \left(\frac{\rho - \rho_0}{\rho_0} \right)^2}{0.37 \left(\frac{\rho - \rho_0}{\rho_0} \right) + 1} \right) \frac{\wp_E(E_f, \beta, q)}{\wp_E(E_f, \beta, 1)} A. \quad (12.349)$$

The dependence of $\wp_E(E_f, \beta, q)$ in Eq. 12.349 on the parameter q is defined by Eq. 12.187 and is presented in Fig. 12.68. In Fig. 12.80, we give

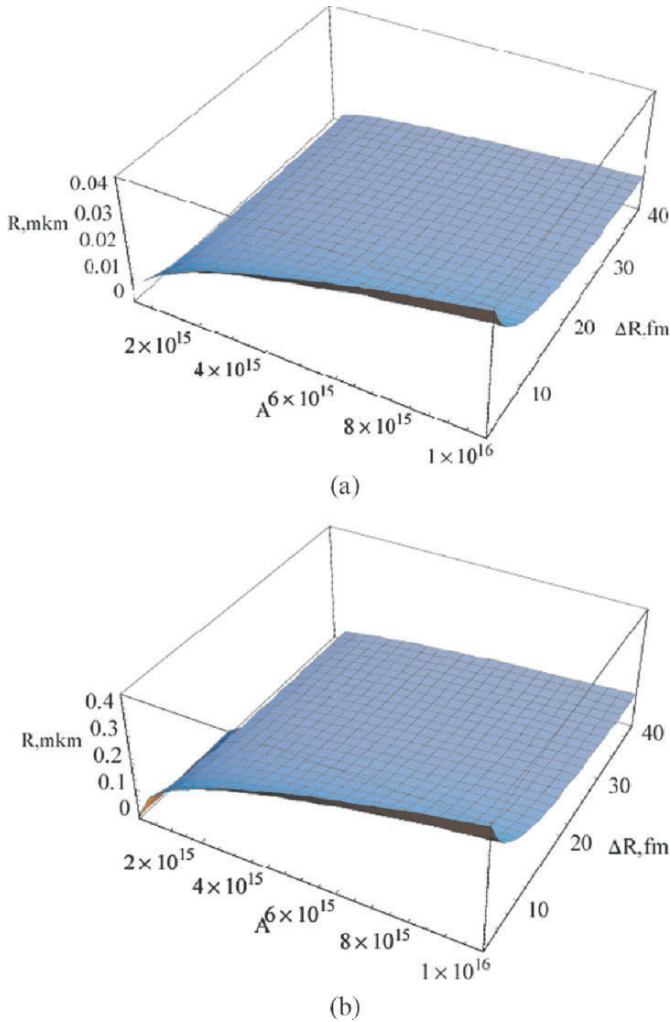


Fig. 12.78. Radius of a cluster versus the shell thickness and the number of nucleons for: standard density of the nuclear substance ρ_0 , (a), density of the nuclear substance equal to $0.01\rho_0$ (b).

the volume part of the strong interaction energy versus the density of the nuclear substance without correlations ($q = 1$). The minimum of the energy corresponds to the density $\rho = \rho_0 = 0.14$.

Figures 12.81 and 12.82 show the dependence of the binding energy on the parameter of nonextensiveness which is related to the degree of correlations between hadrons.

In Fig. 12.81, we display the dependence of the binding energy on the parameter of nonextensiveness at the equilibrium density, and Fig. 12.82

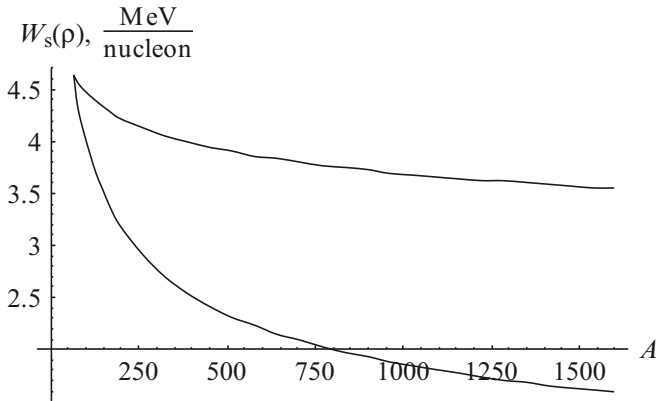


Fig. 12.79. Contribution of the surface to the binding energy of a nucleus versus the mass number with regard to the fractality of the surface.

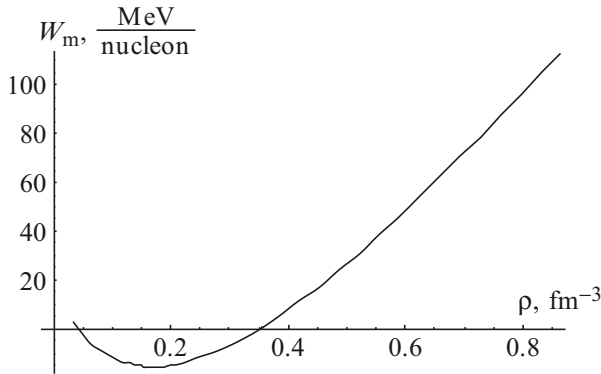


Fig. 12.80. Contribution of the strong interaction to the binding energy of a nucleus (in MeV per nucleon) versus the substance density without regard for correlations ($q = 1$).

shows the surface representing the binding energy of the nuclear substance as a function of both the density and the parameter of nonextensiveness.

Let us estimate the dependence of the parameter of nonextensiveness on the degree of an external action on the system. It follows from the analysis of the asymptotics of the solutions of various kinetic equations with sources and sinks that the main quantity defining both the degree of deviations of the system from the equilibrium and correlations in the system is the flow of energy or the flow of particles. In many cases, the exponent of the asymptotics is related to the intensities of these flows in the phase space. If the energy flow is absent, the solutions pass to equilibrium ones, whereas a growth of the flow intensity is accompanied by the increase in a deviation from the equilibrium.

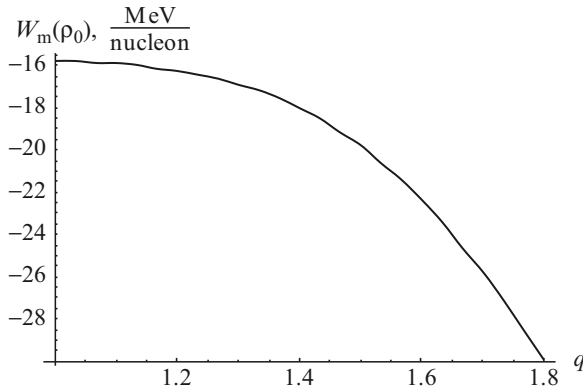


Fig. 12.81. Influence of the correlations on the binding energy of the nuclear substance (in MeV per nucleon) for the equilibrium density of the substance.

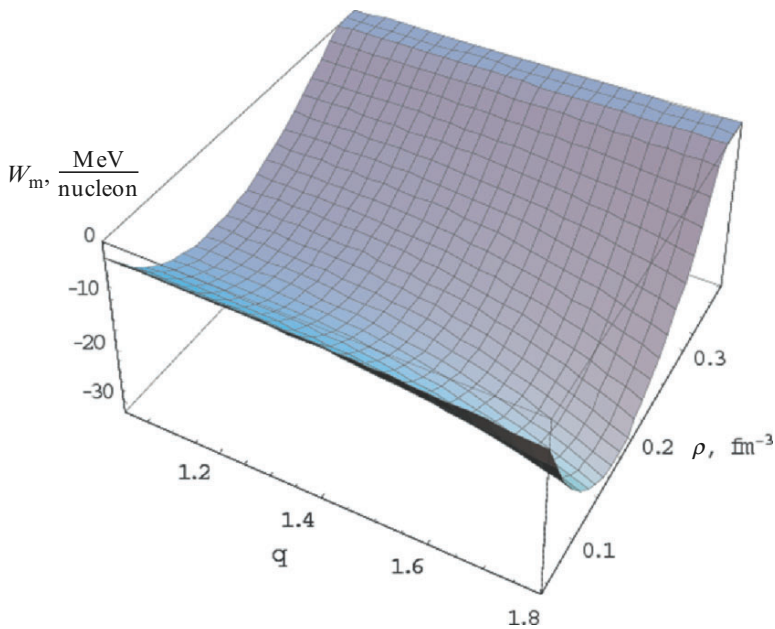


Fig. 12.82. Surface representing the dependence of the binding energy of a nucleus (in MeV per nucleon) on the substance density and the degree of nonextensiveness (correlation).

This dependence can be functionally represented in the form

$$q = \sqrt{1 + \alpha P/P^*}, \quad (12.350)$$

where P^* is the value of the characteristic flow of energy into the system, and α is a free (indefinite) parameter.

One more contribution to the binding energy, which is related to the so-called symmetry-related energy W_β arising by virtue of the Pauli principle, reflects the inclination of the system to pass in the state with the approximately equal numbers of neutrons and protons. Due to the charge symmetry of nuclear forces, the energy at a given density should be an even function of the parameter

$$\alpha_s = N - Z/N + Z = 1 - 2Z/A. \quad (12.351)$$

In the Weizsäcker formula, the symmetry-related energy has the form $W_\beta(A, Z) = c_3 \alpha_s^2 A$ (see Eq. 12.82). It was noted in Ref. 379 that the symmetry-related energy must include a correction connected with the surface energy.

Since a macroscopic nuclear cluster, which arises as a result of the evolution of the above-described nonlinear dynamical processes, has fractal structure of the surface, we use formula Eq. 12.351 for the surface energy and write the expression for the symmetry-related energy as

$$W_\beta(A, Z, D_f) = \alpha_s^2 A \left(c_3 - c_4 \sqrt{c_2} / \sqrt{W_s(A, D_f)} \right). \quad (12.352)$$

It should be mentioned that work Ref. 380 substantiates the ratio between the coefficients $c_4/c_3 = c_2/c_1 = 1.14$ at $c_3 = 30$.

The change of the functional dependence of the symmetry-related energy on the mass number A leads to the significant variations in the stability of nuclear systems with great mass numbers.

Contribution of the Coulomb Energy to the Energy of a Nuclear Cluster. The most complicated situation arises upon the estimation of the contribution of the Coulomb interaction to the binding energy of a macroscopic nuclear structure. Under the assumption of a uniform distribution of charge inside a nucleus, this quantity is presented by the formula $\frac{3}{5} \frac{Z^2 e^2}{R}$. Five-sixths of this contribution of the Coulomb field to the energy of a nucleus consists of the energy of the field outside the nucleus. It is clear that the changes in the environment (especially if the environment is a non-equilibrium dense plasma) should significantly affect the binding energy of the nucleus by inducing a change in that part of the field of the nucleus which exists in plasma outside the nucleus.

As for the part of the Coulomb field which is located in the nucleus, it can be also strongly modified with regard to the creation of pairs of particles of different types from vacuum in the field of the nucleus at sufficiently large Z .

Consider the Coulomb energy of a nuclear cluster in the form of a spherical layer with a fractal structure. The distribution of nucleons over the radius in the cluster defines the dependence of the Coulomb field intensity $E(r)$ on the radius. We estimate the energy of the Coulomb field in the cluster by integrating the field energy density $\frac{E(r)^2}{8\pi}$ with respect to the radius from the internal surface of the cluster to its external surface:

$$W_{\delta in}(A, Z, \delta, \rho) = \frac{Z^2 e^2}{10} \frac{R_\delta(A, \rho, \delta)^5 - (R_\delta(A, \rho, \delta) - \delta)^5}{(R_\delta(A, \rho, \delta)^3 - (R_\delta(A, \rho, \delta) - \delta)^3)^2}. \quad (12.353)$$

The Coulomb field energy outside the cluster can be calculated with regard to the screening in a dense plasma. For the Debye wave number $\kappa_D = \frac{1}{D_s}$ (see Eq. 12.94), we have the relation (see Ref. 272) $\kappa_D r_e = 0.1718 + 0.9283 r_s + 1.591 \cdot 10^2 r_s^2 - 3.8 \cdot 10^3 r_s^3 + 3.706 \cdot 10^4 r_s^4 - 1.307 \cdot 10^5 r_s^5$, where $r_e = (3/4\pi n_e)^{1/3}$, $r_s = m_e e^2 / \hbar^2 r_e$. With regard to the mean parameters of HD, we may approximately write $\kappa_D(n_e) \approx 0.17 (4\pi/3n_e)^{1/3}$. Hence,

$$U_{out}(r, n_e) = (Z e/r) \exp(-\kappa_D(n_e)r), \quad r > R. \quad (12.354)$$

In this case, the energy density of the Coulomb field outside the cluster is $\frac{Z^2 e^2}{8\pi} \frac{\exp(-2\kappa_D r)}{r^2} \left(\frac{1}{r^2} + \frac{2\kappa_D}{r} + \kappa_D^2 \right)$. Calculating the integral of the energy density over the whole volume outside the thin spherical layer of the cluster with radius $R_\delta(A, \rho, \delta)$ and thickness δ , we get

$$W_{\delta out}(A, Z, \delta, \rho, n_e) = \frac{Z^2 e^2}{4} \frac{\exp(-2\kappa_D(n_e)R_\delta(A, \rho, \delta))}{R_\delta(A, \rho, \delta)} (2 + \kappa_D(n_e)R_\delta(A, \rho, \delta)). \quad (12.355)$$

The Coulomb contribution to the binding energy versus the shell thickness and the mass number is shown in Fig. 12.83.

We note that the contribution of the Coulomb field to the binding energy of a nuclear cluster has a more complicated character, than it is presented above. This is related to that the nuclear cluster grows due to the development of a plasma-field structure which has a multilayer character along the radius (see Sec. 12.2.6 and Figs. 12.63–12.65).

The number of layers of this structure $N_l(q)$ is defined by the flow of energy in the system and, hence, upon the theoretical description, by the parameter q .

Thus, the Coulomb energy of a nuclear fractal cluster is the energy of a multilayer condenser with layers possessing a fractal structure.

That is, the Coulomb contribution to the energy of a nuclear cluster consists of the Coulomb energy of layers $W_{\delta in}(A, Z, \delta, \rho)$, Coulomb energy

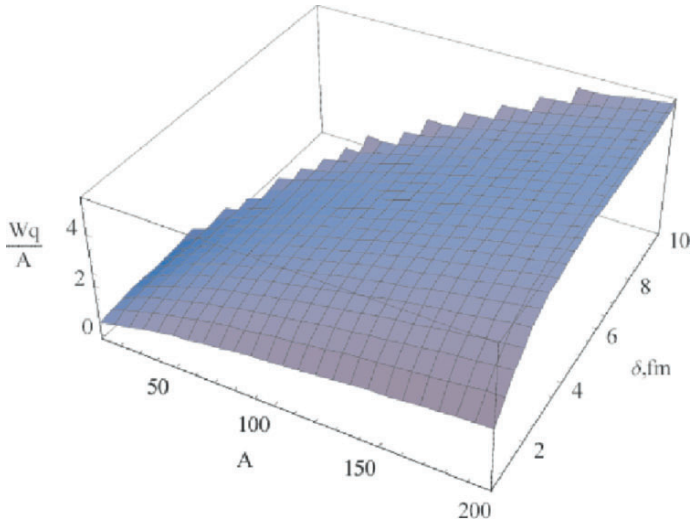


Fig. 12.83. Coulomb contribution to the binding energy versus the shell thickness and mass number.

outside the cluster $W_{\delta out}(A, Z, \delta, \rho, n_e)$, and Coulomb energy of the fields between layers $4\pi\delta N(q) R_\delta^2 \frac{E_{av}^2}{8\pi}$:

$$W_q(A, Z, q, n_e, E_{av}) = 4\pi\delta N(q) R_\delta^2 E_{av}^2 / 8\pi + W_{\delta in}(A, Z, \delta, \rho) + W_{\delta out}(A, Z, \delta, \rho, n_e). \quad (12.356)$$

As a result, the contribution of the Coulomb field to the energy of a nuclear cluster turns out to be dependent practically on all the parameters of the nuclear system.

Taking the above-presented contributions into account, we get the following expression for the binding energy of a nuclear cluster:

$$B(A, Z, q, n_e, E_{av}, D_f) = W_v(\rho, \beta, q) - W_s(A, D_f, \rho, \delta) - W_q(A, Z, q, n_e, E_{av}) - W_\beta(A, Z, D_f). \quad (12.357)$$

As seen, this formula is strongly modified as compared to the ordinary one, which influences, first of all, the conditions of stability of a nuclear cluster. By using the conditions of stability of the nuclear system relative to β - processes

$$\left(\frac{\partial B}{\partial Z}\right)_{n,A} = 0, \quad \left(\frac{\partial B}{\partial A}\right)_{n,Z} = 0, \quad (12.358)$$

we arrive at the relation between the mass number and the charge of a nucleus $Z = Z_\beta(A)$ under condition of its stability. It is natural that this dependence varies upon a change of the parameters of a cluster and can be

different for different types of nuclear structures listed at the beginning of this section.

First of all, consider the binding energy of nuclear structures which are nuclei in the form of liquid drops and bubbles (with the ordinary and fractal dependences of the density on the radius).

In Fig. 12.84, we show the stability line for the standard model of a liquid drop and with regard to the pion condensation.

In Fig. 12.85, we show the dependence of the binding energy on the mass number on the surface of stability drawn in Fig. 12.84.

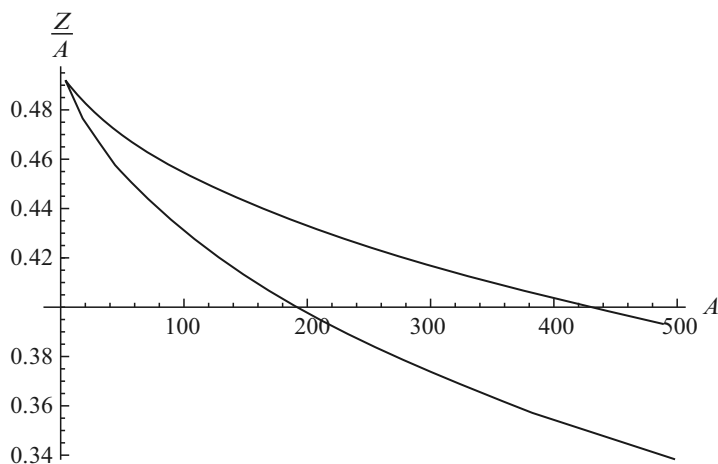


Fig. 12.84. Dependence $Z = Z_\beta(A)$ ensuring the β -stability (see condition Eq. 12.358). The lower curve corresponds to the standard model of a liquid drop, and the upper one takes into account the process of pion condensation.

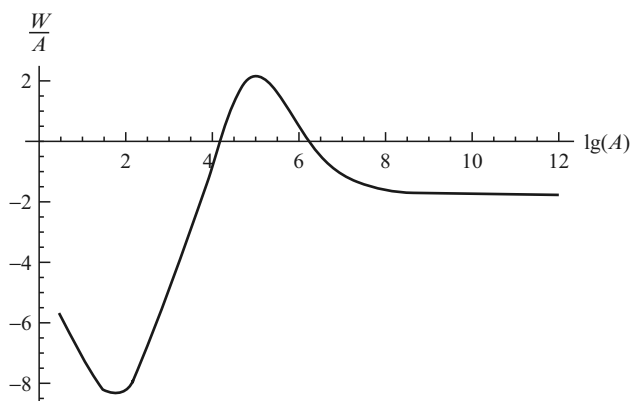


Fig. 12.85. Dependence of the binding energy on the mass number on the surface of stability (see Fig. 12.84).

12.3.4. Active Phase of the Evolution of a Nuclear Cluster in the Form of a System of Shells and the Peculiarities of its Dynamics

We have shown that, at the expense of the dynamical processes running in the system, there arises a *macroscopic nuclear cluster*, which has a complicated structure and a very small radius, i.e., a cluster, being a system of concentric shells with a fractal internal structure.

It is considered usually upon the analysis of the reactions of synthesis that the positive energy yield of the reaction can be only upon the fusion of light nuclei. The reactions with heavy nuclei are energy-gained only in the case of the reactions of decay and are energy-consuming for the reactions of synthesis. However, the situation in our case is not so obvious. Since the cluster has a fractal structure, its surface can even exceed the total surface of fused nuclei. The volume component is not changed in the reactions, and the Coulomb component is strongly suppressed and gives a comparatively small contribution.

Thus, almost all difference in the energies of the final and initial states is accumulated in the surface energy of the cluster.

Consider more comprehensively the binding energy of the nuclear shell.

Binding Energy of a Macroscopic Nuclear Shell. The shell energy has a number of peculiarities, which forces us to consider them separately. These are, first of all:

- macroscopic size of the shell
- contribution of the kinetic energy of electrons, which together with nucleons belong to the shell, to the binding energy
- contribution of the energy of the magnetic interaction of nucleons of the shell with the magnetic field of a current passing in the system to the binding energy

Consider the contribution of the energy of degenerate electrons to the binding energy. The main equations defining a potential distribution in the plasma-field structure, which represents a fine structure of the electron-nucleus shell, are given in the previous section. It has been shown there that the hydrodynamic stage of the evolution of the shell induces the reduction of spatial scales and the appearance of very thin layers of nuclei which are enveloped by much wider layers of the degenerate electron liquid.

In the previous section, we have considered the formation of a structure in an open system within the classical approach. We now show that

the result is preserved in the framework of quantum equations, and the structure constructed from shells appears as well. Consider the final stage of the development of an electron-nucleus structure in the framework of the Thomas–Fermi relativistic equations

$$\frac{\partial}{\partial \rho} \left((R - \rho)^2 \frac{\partial V}{\partial \rho} \right) = -4\pi \frac{e^2}{\hbar c} (R - \rho)^2 (n_e(V(\rho)) - Z n_p(V(\rho))), \quad (12.359)$$

with boundary conditions corresponding to the formation of the structure, rather than with the ordinary conditions describing a single nucleus in vacuum.

Upon consideration of the distribution of electrons in a vacuum shell (i.e., electrons appeared in a nucleus due to the processes of formation of pairs from vacuum in the field a nucleus), it is usually assumed that the distribution of protons is uniform. Upon the consideration of models of the Wigner–Seitz type, the distribution of protons is assumed to be singular (see, e.g., Sec. 12.2.). In our case, after the formation of a macroscopic nuclear cluster in the form of a shell, the distribution of nuclei in each separate layer of the structure is well approximated by a thin spherical shell (see Fig. 12.65) with the thickness $2\sigma_p$:

$$n_p(r) = (n_{p0}/\sqrt{2\pi}) e^{-(r-R)^2/2\sigma_p^2}. \quad (12.360)$$

In the first approximation, the density of nuclei inside a separate layer of the cluster $-\sigma_p < \rho < \sigma_p$ (due to its small thickness) can be considered as constant. Outside the shell (more exactly, between the adjacent layers of the nuclear structure), it can be practically taken as zero.

In the framework of the Dirac equation (see Sec. 12.2.), the system of equations for positrons and electrons is reduced to one second-order equation for the wave function of electrons by the change of variables. As in the previous section, we take into account that $\rho \ll R$ and get a relativistic generalization of the Thomas–Fermi equations in the above-mentioned approximation:

in the region $-\sigma_p < \rho < \sigma_p$ of a nuclear layer,

$$\frac{\partial}{\partial \rho} \left(\frac{\partial \psi}{\partial \rho} \right) = 4\pi \frac{e^2}{\hbar c} \left[\frac{1}{3\pi^2} \left((\omega + \psi(\rho))^2 - 1 \right)^{3/2} - Z n_p \right]; \quad (12.361)$$

and, in the region between nuclear layers,

$$\frac{\partial}{\partial \rho} \left(\frac{\partial \psi}{\partial \rho} \right) = 4\pi \frac{e^2}{\hbar c} \left[\frac{1}{3\pi^2} \left((\omega + \psi(\rho))^2 - 1 \right)^{3/2} \right]. \quad (12.362)$$

We recall that, in these equations, $\psi = e\varphi/m_e c^2$ and $\omega = \varepsilon/m_e c^2$. Each of these equations can be exactly integrated once and is reduced to the form Eq. 12.281.

Let us begin from the region of a nuclear layer. By integrating with regard to the boundary condition $(\partial\psi/\partial\rho)|_{\rho=0} = 0$ at the layer center (at $\rho = 0$, $\psi = \psi_{\max}$), we transform Eq. 12.362 into

$$\left(\frac{\partial\psi}{\partial\rho}\right) = \pm \sqrt{\frac{2}{3\pi} \frac{e^2}{\hbar c} [\psi_u(\psi_{\max}) - \psi_u(\psi)]}, \quad (12.363)$$

$$\begin{aligned} \Psi_u(\psi) = & \frac{1}{8} (\omega + \psi) \sqrt{(\omega + \psi)^2 - 1} \left[2(\omega + \psi)^2 - 5 \right] \\ & + \frac{3}{8} \ln \left((\omega + \psi) + \sqrt{(\omega + \psi)^2 - 1} \right) - \beta\psi. \end{aligned} \quad (12.364)$$

In the regions between nuclear layers, the equation for the potential acquires the same form, but with the function $\Psi_u(\psi)$ without the last term in Eq. 12.364.

These equations define the trajectories of the system on the phase plane. The plots of trajectories of the system on the phase plane for various values of the parameter β are given in Fig. 12.86 (a).

The initial point corresponding to the maximum of the potential $U(\rho_0)$ defines the right boundary of the interval of potentials, where the shell exists.

The form of trajectories in the region, where the density of nuclei is very small, is qualitatively different from that of the above-presented ones by that the phase curves in this region are monotonous and have no second point with the zero derivative at the edge (see Fig. 12.86 [b]).

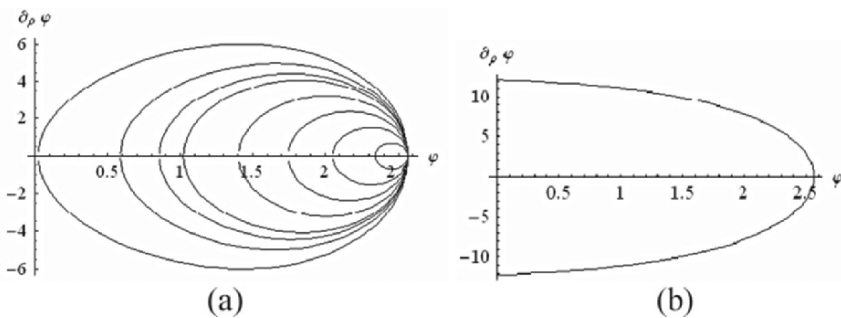


Fig. 12.86. Trajectories of the solutions of the relativistic Thomas-Fermi equation for various values of the parameter β : inside a nuclear cluster (a); between the nuclear layers of the cluster (b).

Two above-presented families of curves intersect at that point of the phase space which corresponds to the transition from the layer region with the dominance of the positive charge (nuclei) to the region with the dominance of the negative charge of electrons. As a result, we get the full family of trajectories.

The solution of the Thomas–Fermi equations in the whole region can be also obtained numerically

The potential, which is constructed by these solutions, in line with Figs. 12.60–12.62.

We now pass to the description of the dynamics of a nuclear cluster-shell.

Dynamics of a Macroscopic Nuclear Cluster-Shell. The study of the motion of shock waves in stars within the approach of self-similar motions was begun in the works of L. Sedov and K. Stanyukovich (see Ref. 234). That investigation was based on the equations describing a spherically symmetric motion of a gas with regard to the gravity force (see Ref. 383).

For the motion of the shell in a target, it is necessary to generalize the appropriate equations because the gravity force can be neglected, and the motive force of a shell collapse is the surface energy of external layers and the energy released upon the fragmentation (evaporation) of nuclei from the surface of a shell nuclear cluster.

The dynamics of a shell along the radius is defined by the forces acting on it. They are, first of all, the forces of pressure from the side of the internal region which counteract to the compression of the shell, as well as the external pressure and the surface tension causing its compression.

In the process of evolution of a cluster, it absorbs the flows of nuclei of the target from the side of the internal surface, whereas the external surface of the cluster evaporate fragments with various mass numbers. The distribution of these fragments is defined by Eqs. 12.360, 12.361.

As a result, the energy variation upon a change of the shell radius with time is

$$\Delta W(R, q) = \int_1^{\infty} P_{\mathbf{p}}(A, Z_{\beta}(A)) \varepsilon_{\mathbf{p}}(A, Z_{\beta}(A)) dA. \quad (12.365)$$

Here, we use the dependence between the mass number and the charge on the line of β -stability [$Z = Z_{\beta}(A)$].

We denote the shell thickness as Δ , its external radius as $R(t)$, the fractal dimensionality of the cluster surface as D_f , and the mass inside the shell with radius R and thickness Δ as $M(R, \Delta)$.

Then the effective area of the surface turns out to be equal to $S_{eff} = R^{\gamma D_f}(t)$, where $\gamma < 1$ ($\gamma \approx 0.86$). The energy concentrated near the surface (at the expense of the surface energy) is proportional to the area, and the force acting on the shell is

$$F_S \approx \frac{\partial S_{eff}}{\partial R} \approx R^{\gamma D_f - 1}. \quad (12.366)$$

The acceleration of the shell as a result of the action of the surface tension is

$$g_s(R) \approx \frac{F_S(R)}{M(R, \Delta)} \approx \frac{R^{\gamma D_f - 1}}{4\pi R^2 \rho^* \Delta} \approx \frac{1}{R^{3 - \gamma D_f}}. \quad (12.367)$$

The change of the shell energy $\Delta W(R, q)$ upon a variation of its radius defines the additional force acting on the shell in the active phase of its evolution (in the phase of nuclear combustion):

$$F_{nu} \approx \frac{\partial \Delta W(R, q)}{\partial R}, \quad g_{nu}(R) \approx \frac{F_{nu}(R)}{M(R, \Delta)} = \frac{F_{nu}(R)}{4\pi R^2 \rho^* \Delta}. \quad (12.368)$$

The motion of a self-compacting shell toward the center is similar to the motion of a collapsing star and is described by the equations

$$\frac{\partial \rho}{\partial t} + \frac{\partial(\rho v)}{\partial r} + \frac{2(\rho v)}{r} = 0, \quad (12.369)$$

$$\frac{\partial v}{\partial t} + v \frac{\partial(v)}{\partial r} + \frac{1}{\rho} \frac{\partial P}{\partial r} = -g(r), \quad g(r) = g_s(r) + g_{nu}(r). \quad (12.370)$$

Let us write the condition of the stability of a shell (the zero value of the sum of accelerations at the expense of both the growth of a pressure with decrease in the radius and the action of the surface forces):

$$\frac{1}{\rho} \frac{\partial P}{\partial r} = -g(r). \quad (12.371)$$

We take the equation of state in the approximation of a barotropic medium, $P = P(\rho)$. Then $\frac{dP}{dr} = \frac{dP}{d\rho} \frac{d\rho}{dr}$, and, after the substitution of it in the equation of equilibrium, we get

$$\frac{d\rho}{dr} = -\frac{g_0}{(dP/d\rho)} \frac{r_0^{3 - \gamma D_f}}{r^{3 - \gamma D_f}} \rho(r). \quad (12.372)$$

The derivative $\frac{dP(\rho)}{d\rho}$ for the polytrope $P(\rho) = P_0 \left(\frac{\rho}{\rho_0}\right)^{\gamma_a}$ reads

$$\frac{dP(\rho)}{d\rho} = \gamma_a \frac{P_0}{\rho_0} \left(\frac{\rho}{\rho_0}\right)^{\gamma_a - 1}. \quad (12.373)$$

Then

$$\frac{1}{\rho_0} \frac{d\rho}{dr} = - \frac{g_0}{\left(\gamma_a \frac{P_0}{\rho_0} \left(\frac{\rho}{\rho_0} \right)^{\gamma_a - 2} \right)} \frac{r_0^{3 - \gamma D_f}}{r^{3 - \gamma D_f}}. \quad (12.374)$$

The entire process of evolution of the shell is initiated by an external pulse which defines the parameters of the mode with sharpening and then the parameters of the structure, being simultaneously the parameters of the binding energy of the nuclear cluster. The resulting distribution of fragments by mass upon the evaporation is defined by a quasistationary state.

The resulting motion is a strongly nonlinear process, whose direction depends on many parameters and requires a careful analysis.

Thus, the described process can be represented on the whole as the process of transformation of a part of the rest mass of the initial substance scanned by the shell during its evolution in the target bulk in the energy of the shell (the free and bound ones).

Upon approaching the center, the system of concentric layers reaches its maximum parameters prior to the loss of stability and its fracture. As an important peculiarity of the extreme state of the shell substance, we mention the process of protonization of free neutrons in the shell substance at attained densities which is considered in Chap. 11.1.

The Explosion Phase of a Macroscopic Nuclear Shell. Then there occur the dispersion of fragments with various sizes from the region, where the wave-shell has exploded, and the interaction of products of the explosion of the nuclear megacluster with nuclei of the target substance. During this interaction, there are the running ordinary high-energy nuclear processes which can lead to the formation of a wide spectrum of products of the reactions (including the radioactive products). We recall that the *synthesis of nuclei upon the motion of the wave toward the center occurs as a result of the adiabatic processes of fragmentation and, like the phenomenon of cluster radioactivity, is running in the class of stable nuclei.*

It is obvious that the statistical distribution of fragments of a decaying shell by mass numbers is the reflection of a nonequilibrium state of the macroscopic nuclear structure – the shell.

Analogous processes of fracture of a compound nucleus upon the collision of heavy ions lead to the power distributions of fragments of the form $\approx A^{-2.65}$ by mass numbers. Such mass spectra, which are observed upon the fracture of a compound nucleus upon its great excitation and upon the fracture of a shell as a result of the loss of its stability, can be explained within statistical models with regard to nonextensiveness (see Ref. 280).

In the nonequilibrium case, the description could be carried out in the kinetic approximation. However, it is much more convenient to describe the processes of fragmentation by using the Tsallis nonextensive statistics (see Ref. 289) (see Sec. 12.2).

Let us assume that the distribution of fragments satisfies the generalized canonical probability distribution (see Refs. 288, 287) by the masses A_k and the charges Z_k of fragments:

$$P_{\mathbf{p}}(A_k, Z_k) = (\Omega_q(A_k, Z_k))^{-1} [1 + (q-1)\beta\varepsilon_{\mathbf{p}}(A_k, Z_k)]^{-1/(q-1)}. \quad (12.375)$$

Here,

$$\begin{aligned} \Omega_q(A_k, Z_k) &= \sum_k (1 + (q-1)\beta\varepsilon_k(A_k, Z_k))^{-1/(q-1)}, \\ \varepsilon_k(A_k, Z_k) &= T_k(A_k, Z_k) + B(A_k, Z_k, q, n_e, E_{av}, D_f), r \\ T_k &= \frac{p_k^2}{2m_k} = \frac{p_k^2}{2(m_n(A_k - Z_k) + m_p Z_k)}, \end{aligned} \quad (12.376)$$

and B is the binding energy of nuclei (see Eq. 12.358). The derived distribution has a power asymptotics.

The main characteristics affecting both these processes and the exponent of the asymptotics of the distribution by mass numbers are the internal correlations existing in the nuclear macrostructure. We may consider nuclides as “words” of a dynamical system, from which they are evaporated. The experimental study of the statistical composition of nuclei created as a result of the decay of the shell is, from this viewpoint, the investigation of the statistical properties of the “language” of a dynamical system (see Refs. 384–387).

In Fig. 12.87, we show the Zipf distribution (the frequency distribution of “words”—nuclei) for the earth crust and the sun system.

It is seen from the figure that the distribution for the Earth crust has form typical of a Markov process, whereas the distribution for the sun system has form typical of a dynamical process with high correlations.

In Fig. 12.88, we demonstrate the distribution functions for evaporated nuclei of a nuclear cluster-shell created in the Kiev experiments.

It follows from the analysis of the plots that the distribution functions for nuclei derived in the experiments are close to their abundance in the sun system and occupy some region between the distributions in the sun system and the earth crust (Figs. 12.89, 12.90).

The difference of these distributions is especially clearly seen upon their representation on the double logarithmic scale (see Figs. 12.91, 12.92).

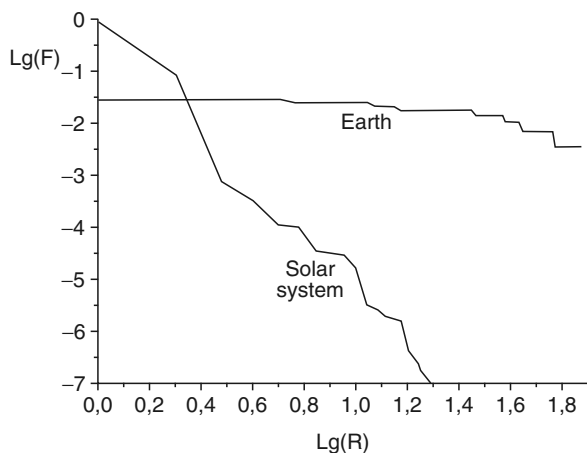


Fig. 12.87. Zipf distribution for the frequencies of nuclei in the Earth crust and the sun system on the double logarithmic scale.

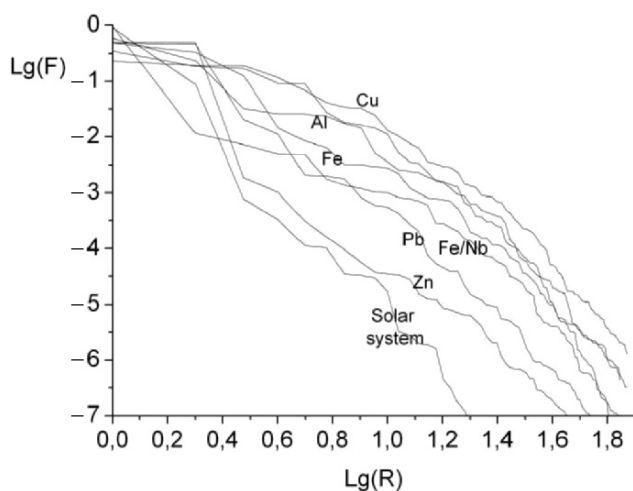


Fig. 12.88. Distribution functions for nuclei evaporated from a nuclear shell in the Kiev experiments together with the distribution for the sun system.

The distribution function of nuclei, being products of the nuclear regeneration of exploded targets in the experiments performed at the Electrodynamics Laboratory “Proton-21”, has form typical of non-Markov systems with strong correlations, whereas the distribution function of the natural elements is close to that for Markov processes.

A power asymptotics of this distribution is obvious:

$$p(R) = A/(B + R)^\xi, \quad (12.377)$$

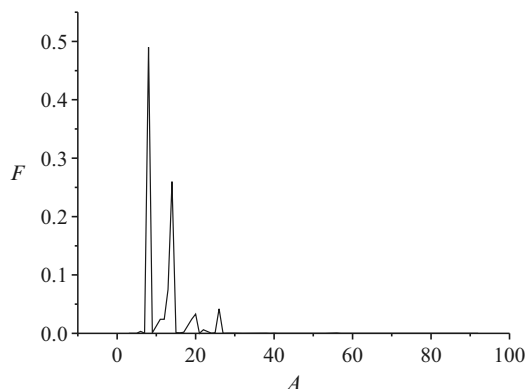


Fig. 12.89. Abundance of elements in the earth crust.

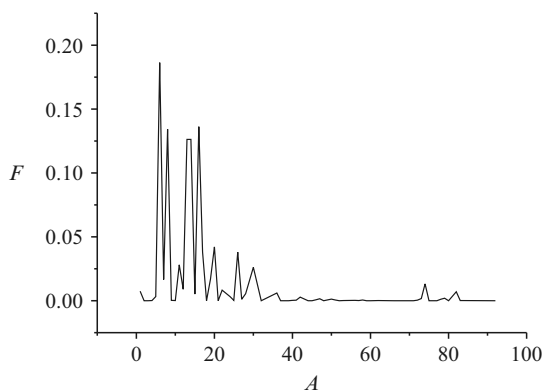


Fig. 12.90. Abundance of elements formed in the Kiev experiments.

where the exponent $\xi \approx 2.786$, which corresponds to the correlation decay degree

$$\alpha = \frac{1 + \xi}{2}. \quad (12.378)$$

In the analysis of the experimental data, it is convenient to construct the frequency dictionary of a language of the dynamical system (See Refs. 186, 187) representing our experiments.

The Zipf frequency distribution for the abundance of elements is given in Fig. 12.93.

The comparison of the results derived at the Electrodynamics Laboratory “Proton-21” with those following from thermodynamical considerations allows us to conclude that the fractal dimensionality of the surface of a nuclear cluster created in the experiments $D_f = 1.8$.

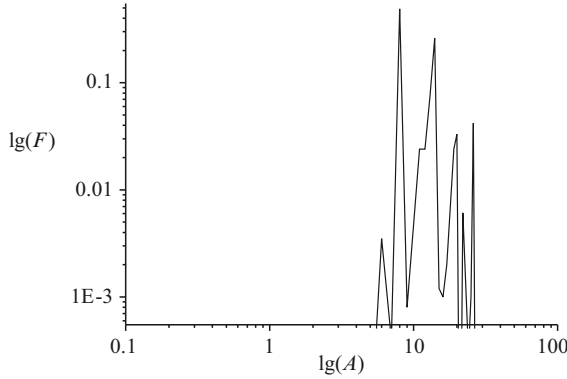


Fig. 12.91. Abundance of elements in the earth crust (on the double logarithmic scale).

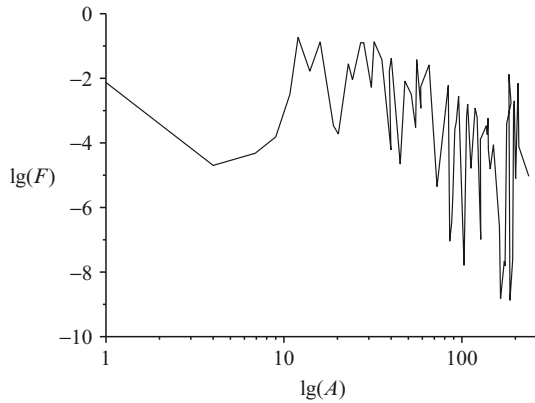


Fig. 12.92. Abundance of elements formed in the Kiev experiments (on the double logarithmic scale).

These results agree with the above-presented estimates of the fractal dimensionality and the parameter of nonextensiveness. The imbalance of the system and the existence of correlations between the states of particles in the system lead also to the nonextensiveness of the quasistationary distributions of particles by energy.

It has been shown in Sec. 12.3.3 that such distributions with power asymptotics at high energies are the solutions of kinetic equations and correspond to the correlations in the system which appear as a result of the presence of the flows of energy or particles. The power asymptotics of the distribution of particles by energy are revealed in the experiments in all parts of the spectrum accessible for measurements. The exponents in the distribution functions as functions of energy are approximately the same in

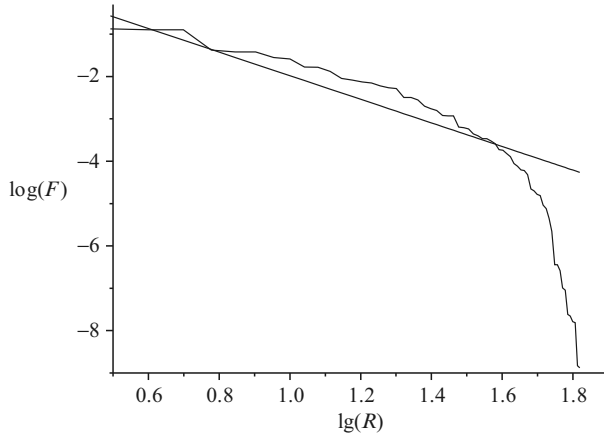


Fig. 12.93. Frequency distribution of nuclei formed in the experiments.

all the regions (see Fig. 4.7), $f(E) \propto E^{-2.7}$, and correspond to the above-considered states with strong correlations.

12.4. Scenario of the Development of a Self-Organizing Nucleosynthesis in the Estimates by the Physical Models Presented in this Work

The temporal dependences of pulses of the voltage, current, and power of a beam, for which the estimates and calculations were performed (see Figs. 12.94–12.96).

The decrease in the surface area where the beam interacts with a target (in the estimates, we considered a Cu target) as a result of both the beam focusing and the plasma dynamics is presented in Fig. 12.97.

The characteristic depth, at which the electron beam loses its energy, versus time is shown in Fig. 12.98.

The ratio of a beam current to the critical current near the diode is presented in Fig. 12.99.

The time moment, at which the critical current is exceeded (15 ns), is the start of the mode with sharpening.

The volume of the region where the beam interacts with a target versus time is presented in Fig. 12.100.

After the start of the mode with sharpening, the interaction region volume (the region where the beam energy is absorbed) is sharply reduced.

The temporal dependence of the logarithm of the ratio of the impact parameter of an action to the limit value (the experimental estimate) is presented in Fig. 12.101.

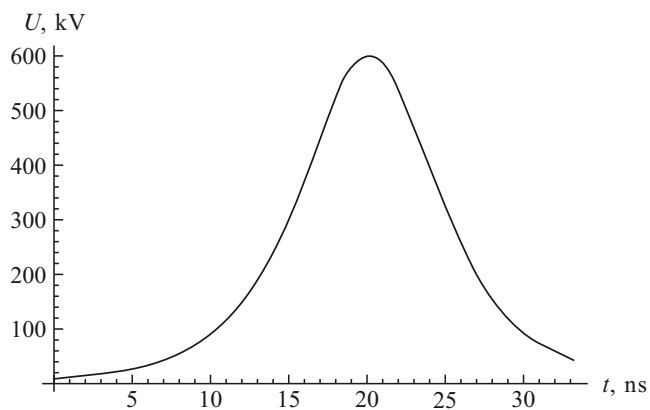


Fig. 12.94. Temporal dependence of pulses of voltage of a beam.

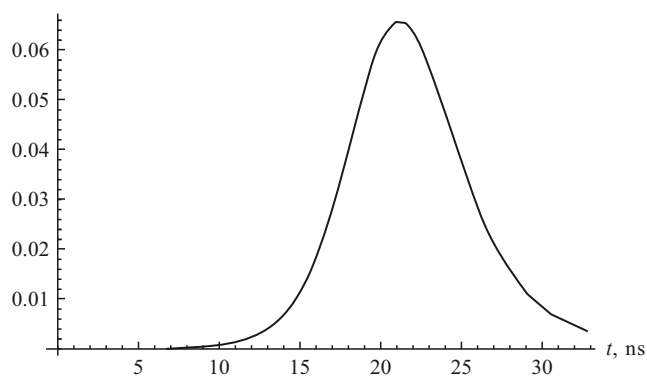


Fig. 12.95. Temporal dependence of pulses of current of a beam.

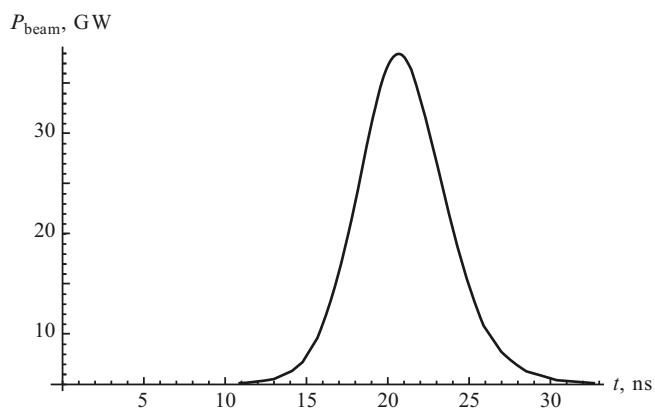


Fig. 12.96. Temporal dependence of pulses of power of a beam.

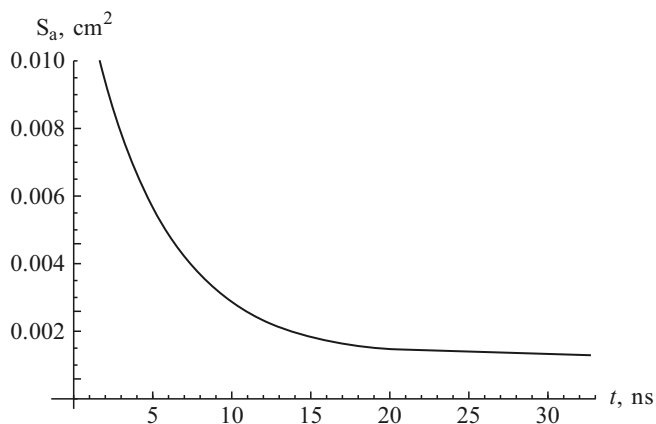


Fig. 12.97. The decrease in the surface area where the beam interacts with a target.

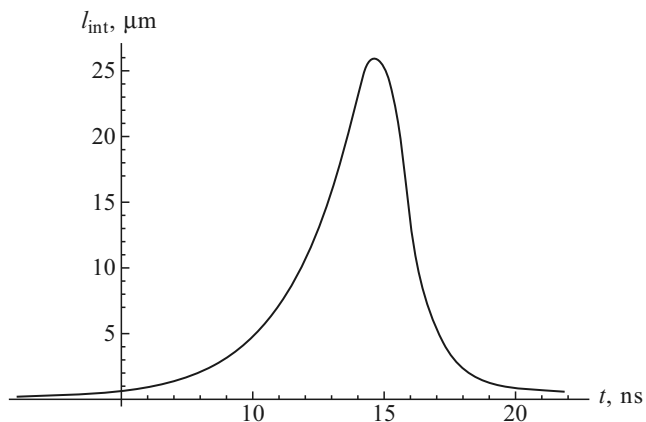


Fig. 12.98. The characteristic depth, at which the electron beam loses its energy, versus time.

It is seen that the excess of the impact parameter above its limit value also happens near the beginning of the mode with sharpening.

The temporal dependence of the estimated pressure in the near-surface layer of a target is presented in Fig. 12.102.

The temporal dependence of the number of nuclei of Cu in the interaction region is presented in Fig. 12.103.

The energy (in keV) per Cu atom in the region of interaction of the beam and the target versus time is presented in Fig. 12.104.

To fully ionize Cu, it is necessary to spend about 45 keV per nucleus, therefore the supplied energy is sufficient for the full ionization of ions of a target in the region of interaction with the beam.

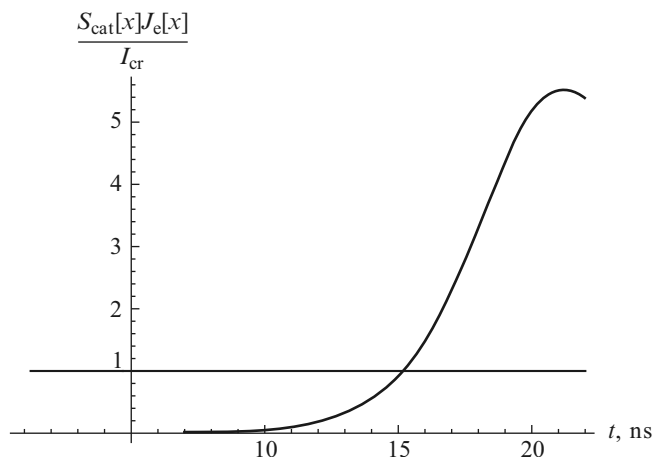


Fig. 12.99. The ratio of a beam current to the critical current near the diode.

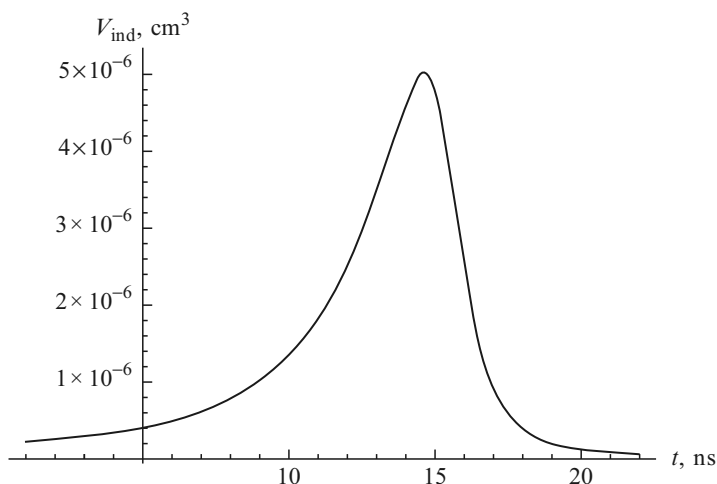


Fig. 12.100. The volume of the region where the beam interacts with a target versus time.

The plot of the distance passed by a nonlinear acoustic wave (in μm) up to its overturn as a function of the time moment when electrons of the beam begin to interact with the surface is presented in Fig. 12.105.

As long as the parameters of a beam do not satisfy the impact criterion, the nonlinear wave can steepen practically only at the center. Beginning from the time moment when the impact criterion becomes to hold, the overturn point approaches the surface. In the one-fluid mode, the maximum increase in the density occurs at depths of 100 to 150 μm , as is indicated by the solutions of hydrodynamic equations.

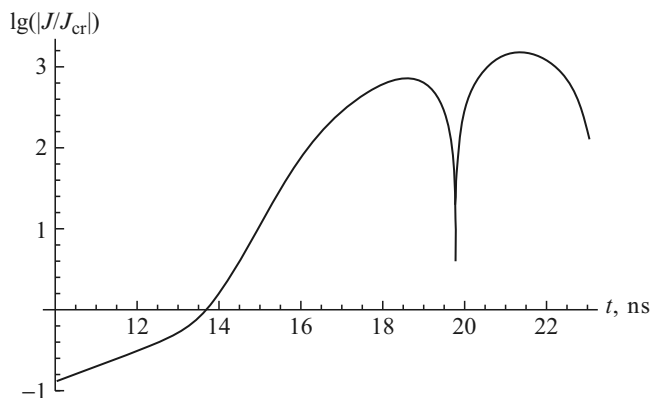


Fig. 12.101. The temporal dependence of the logarithm of the ratio of the impact parameter of an action to the limit value (the experimental estimate).

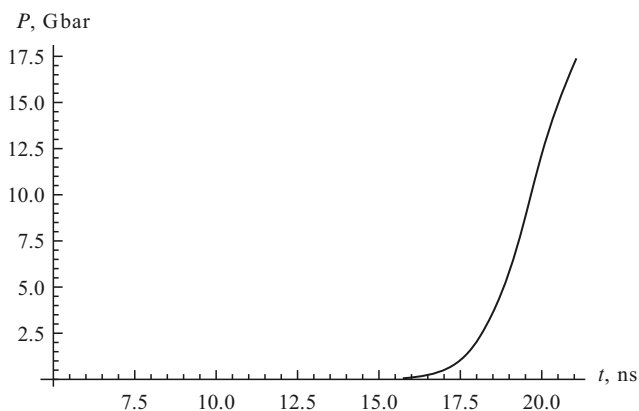


Fig. 12.102. The temporal dependence of the estimated pressure in the near-surface layer of a target.

As shown in Sec. 12.2., the development of an instability leads to the appearance of plasma-field structures and the reduction of their scale up to extremely small values $\leq 10^{-10}$ cm.

The special measurements of X-ray emission in the region of low energies showed a sharp decrease in the emission intensity and its absence at energies below 10 keV. Such effects are usually related to the interaction of the emission with a dense plasma and allow one to estimate its density in the region of HD: $n_p \approx 10^{30} \text{ cm}^{-3}$.

Simple geometric relations define the radius of spherical layers R_b , their thickness Δ_b , and the mean density of a substance in a layer ρ_{av} .

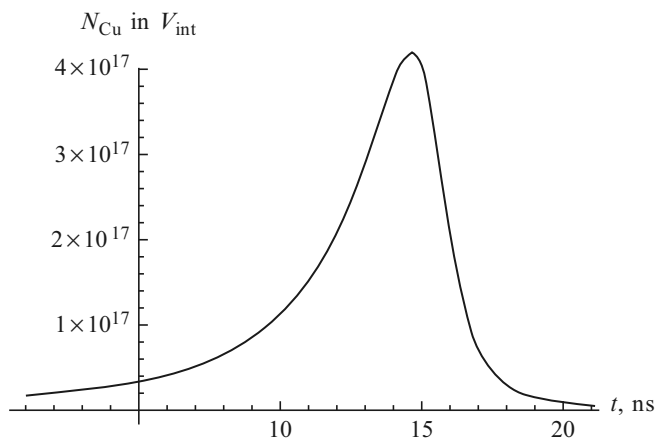


Fig. 12.103. The temporal dependence of the number of nuclei of Cu in the interaction region.

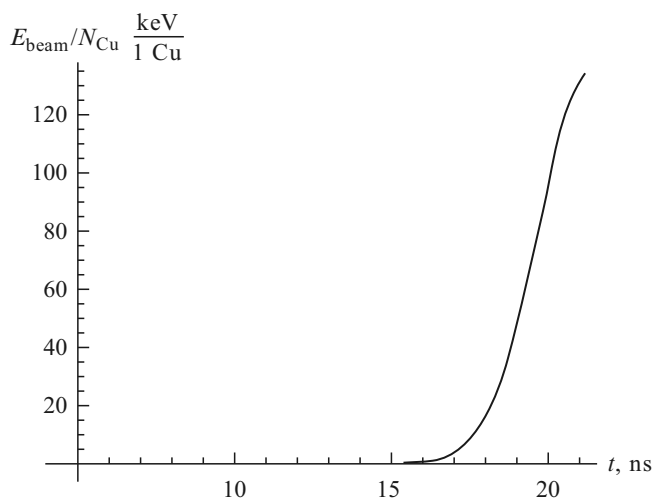


Fig. 12.104. The energy (in keV) per Cu atom in the region of interaction of the beam and the target versus time.

The above-presented estimates yield that about $A_{b0} \approx 10^{17}$ nucleons are present in a thin spherical layer (a shell). The dependence of the external radius of the shell with thickness $\leq 10^{-10}$ cm and with density $n_p \approx 10^{30} \text{ cm}^{-3}$ on the number of nucleons in it is presented in Fig. 12.106.

The shell radius corresponds to that the mentioned critical parameters are established in the region where the extremely nonlinear effects are realized upon the fulfillment of the impact criterion (see above).

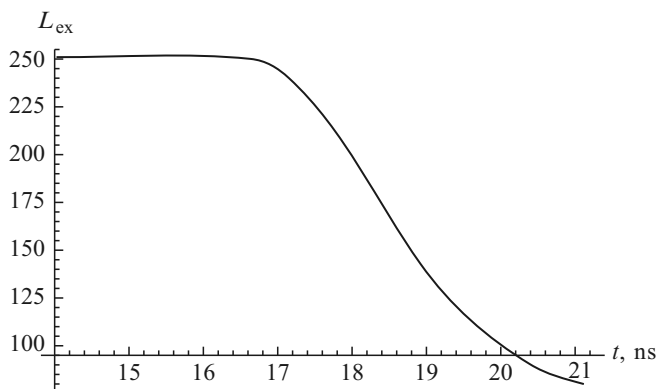


Fig. 12.105. The plot of the distance passed by a nonlinear acoustic wave (in μm) up to its overturn as a function of the time moment.

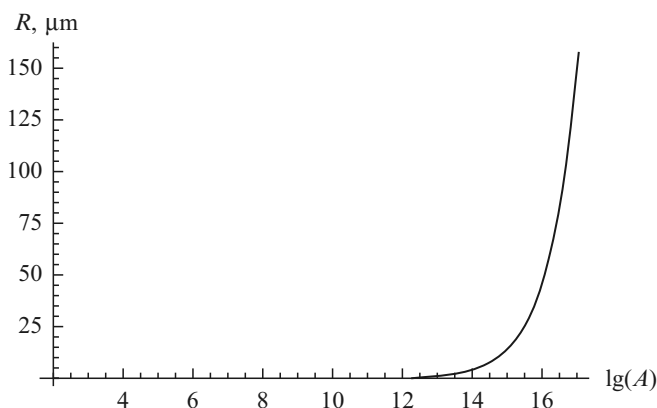


Fig. 12.106. The dependence of the external radius of the shell with thickness $\leq 10^{-10}$ cm and with density $n_p \approx 10^{30} \text{ cm}^{-3}$ on the number of nucleons in it.

Upon the attainment of the indicated parameters, there appear the effect of suppression of the Coulomb repulsion of nuclei at the expense of quantum size effects (the effect of “pairing” of nuclei) and the fast growth of a cluster in the shell volume. There arises the “active” phase of the process, namely the “active” phase of the dynamics of the shell. To estimate this phase, we will use the model of a fluid drop for nuclei (the Weizsäcker model) and will describe the shell within a modified Weizsäcker model which is the simplest model for a fractal nuclear macrobubble-shell.

The initial ratio of protons and neutrons is set by a target substance, and the number of nucleons is defined by the dynamics of the beam and the target and by a value of the interaction volume with regard to the above-presented estimates: A_{b0} is of the order of 10^{17} . In this case, the

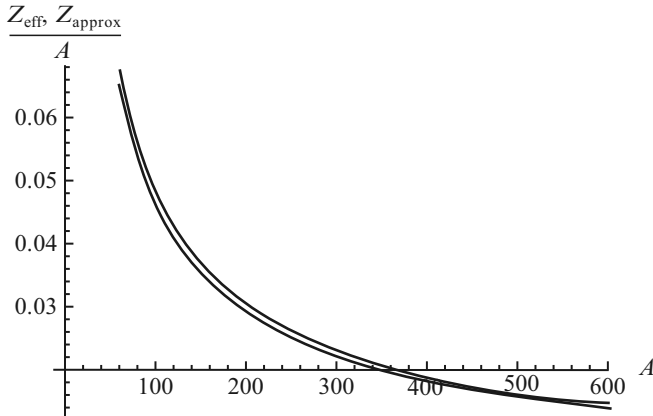


Fig. 12.107. A change of the effective charge.

contribution of the Coulomb interaction is given by the Migdal model of condensation. In this model, as a result of the processes of condensation of mesons and electrons, the effective charge of a nucleus decreases sharply at great distances from it and is expressed by the relation which is well approximated by the function $Z_{\text{eff}} \approx A^{1/3}$. A change of the effective charge is shown in Fig. 12.107.

The contribution of the surface energy can be found in view of the fractality of the structure:

- For a continuous cluster, the area of the surface is proportional to that of the external surface, i.e., to the square of the radius and to $A^{2/3}$.
- For a completely disjointed cluster, its surface is proportional to A , since it is the sum of the area of all nuclei.

In the general case, the exponent as a function of A is in the interval between $2/3$ and 1 , namely it equals $2/3 + \Delta\gamma$, where $0 < \Delta\gamma < 1/3$.

Two plots for components of the binding energy per nucleon in the shell are presented in Figs. 12.108–12.109.

The relative contribution of the Coulomb energy decreases rapidly with increase in the mass number, which is seen in the first plot. The contribution of the surface energy depends on the irregularity of a cluster (whose parameter lies, as explained above, in the interval $0 < \Delta\gamma < 1/3$).

The second plot demonstrates the surface contribution versus its irregularity for a mass number A_{b0} .

The binding energy of the shell per nucleon versus the mass number for three values of the fractality is presented in Fig. 12.110. The upper, middle, and lower curves correspond, respectively, to the slight, medium, and maximum fractalities.

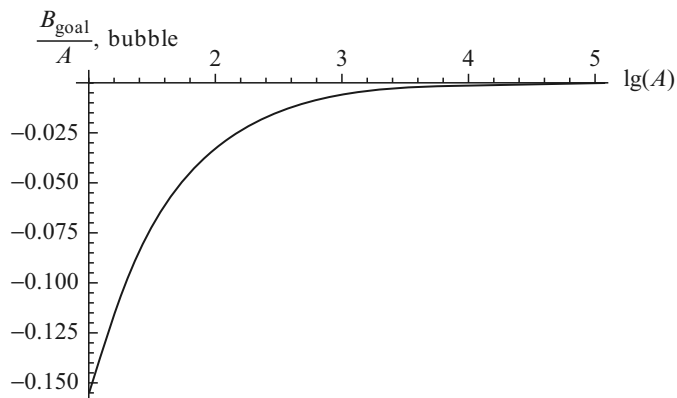


Fig. 12.108. The Coulomb energy per nucleon in the shell.

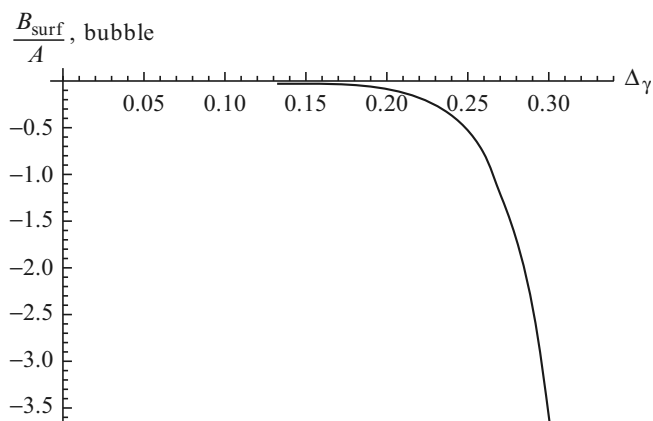


Fig. 12.109. The surface energy per nucleon in the shell.

The dependence on the fractality is seen more clearly in Fig. 12.111, where we give the binding energy per nucleon for a shell with the fixed mass number $A_{b0} \approx 10^{17}$.

Upon the motion of a superdense shell to the center, it should pass about $1.055332 \cdot 10^6$ atomic layers of the target, each having a thickness of about $1.42135 \cdot 10^{-8}$ cm.

While passing each layer, the shell absorbs nucleons, whose number depends on the running radius of a layer $R_b(\text{cm})$: $9.50379 \cdot 10^{17} R_b^2$.

Since the nuclear substance is a noncompressible liquid, its volume does not change in nuclear reactions and, hence, the volume component of the binding energy which is proportional to the number of baryons is constant.

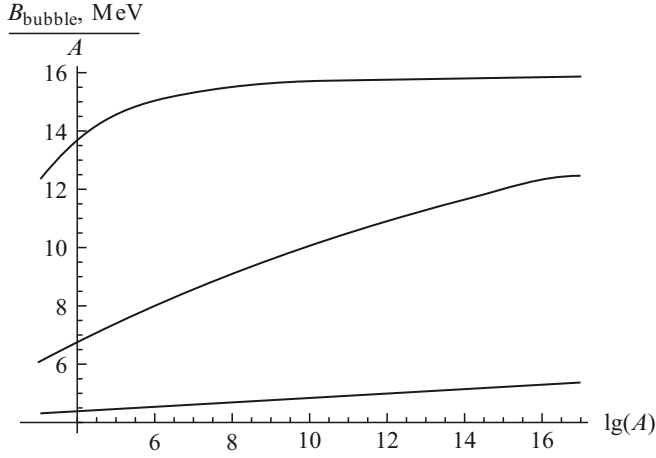


Fig. 12.110. The binding energy of the shell per nucleon versus the mass number for three values of the fractality.

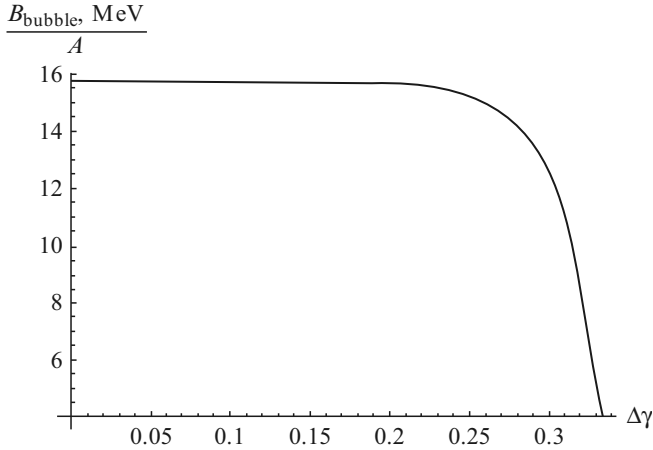


Fig. 12.111. The binding energy per nucleon for a shell with the fixed mass number $A_{b0} \approx 10^{17}$.

In the nuclear cluster, the Coulomb component of the binding energy of nuclei of the target is extremely small and, in the first approximation, can be neglected. Thus, upon the growth of the cluster at the expense of the nucleons entering it, the total Coulomb and surface energies of nuclei of the target from the nearest atomic layer transform into the surface energy of the cluster. Roughly saying, the cluster is a structure similar to a 3D system of web threads. Under assumption that the next layer of nuclei of the target transforms into a web thread, the preservation laws yield that the

length of a new web thread is about $4.80724 \cdot 10^{17} R_b^2$ fm, and its cross-section area is about 14 fm^2 (at each step). Such a cross-section area corresponds to the mass number of elements, from which a “web thread” is built, to be of the order of 4 to 5. That is, we may assume that the “web thread” is built from structures similar to α -particles.

As any macroscopic system, the nuclear macroscopic cluster evaporates fragments from the surface. The evaporated fragments must mainly consist of the most stable structures corresponding to the running state of the shell.

We have shown above that the thermodynamical conditions of equilibrium (without regard for the internal structure of the nuclear substance) lead to the dependence of the mass number of the most stable nucleus on the cluster density which is presented in Fig. 12.112.

As seen, the mass number of the most probable nucleus, being condensed from the nuclear substance, is slightly changed in the huge range of densities and is close to the mass numbers of nuclei from the group of Fe. Though the account of the structure of the nuclear substance can change the ratios of probabilities of condensed structures, we take the simplest model and assume that the evaporated nuclei are those of Fe.

At each step, the evaporated nuclei of Fe are condensed and form an equilibrium substance. In this simple model of a homogeneous nuclear substance, the number of evaporated nuclei of Fe turns out to be about $1.48735 \cdot 10^{16} R_b^2$.

By comparing the number of nucleons which belong to the cluster and are evaporated from its external surface, it is easy to see that the shell

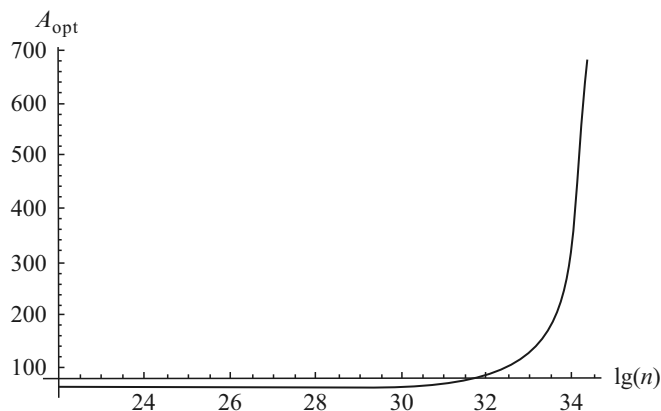


Fig. 12.112. The dependence of the mass number of the most stable nucleus on the cluster density.

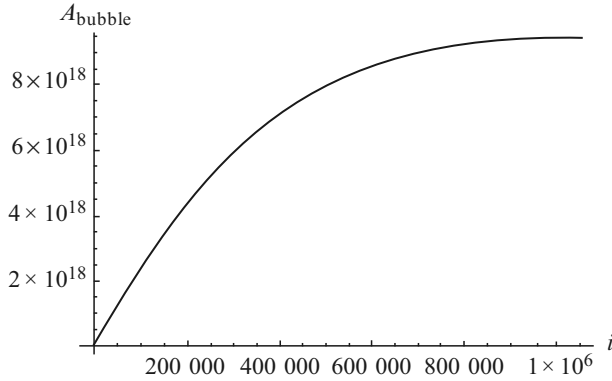


Fig. 12.113. The law of increase of the mass number of the shell as a function of the layer number i .

mass increases while approaching the center. The law of increase of the mass number of the shell as a function of the layer number i ,

$$A_b(i) = 1 \cdot 10^{17} + 7.91011i(3.34117 \cdot 10^{12} - 3.16599 \cdot 10^6 i + i^2)$$

is demonstrated in Fig. 12.113.

During the evolution, the shell absorbs $9.29714 \cdot 10^{18}$ nucleons.

While moving to the center, the shell density increases with the number of a layer as

$$\frac{2.66434 \cdot 10^9(1 \cdot 10^{17} + 1.17463 \cdot 10^{17}(0.015 - 1.42135 \cdot 10^{-8}i)^2)}{(0.015 - 1.42135 \cdot 10^{-8}i)^2}.$$

This growth is especially fast near the center, as is seen from Fig. 12.114.

During its motion to the center, the shell accumulates the energy at the expense of a growth of the cluster surface at each step. As a result, the shell energy at the step j becomes (in MeV)

$$E_{\text{bubble}}(j) = 48.7147j^2(3.34117 \cdot 10^{12} - 3.16599 \cdot 10^6 j + j^2).$$

The dynamics of a growth of the shell energy is shown in Fig. 12.115.

The shell moves to the center under the action of several forces. First, due to its small thickness, there appear the great gradients of density near the surface. The great gradients induce the surface forces compressing (accelerating) the shell. Moreover, the “reactive” force of “evaporated” nuclei acts on the shell in the same direction.

The distribution of the surface energy as a function of the radial coordinate and, hence, the surface force depend on the cluster fractality.

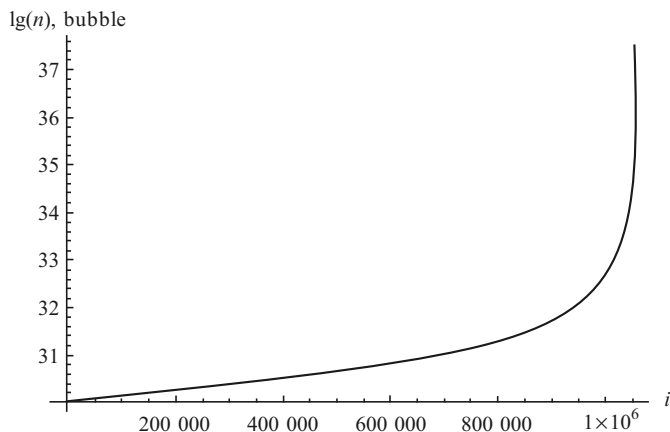


Fig. 12.114. The shell density increase with the number of a layer.

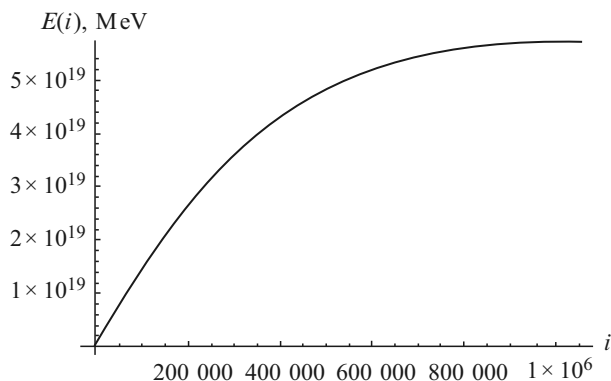


Fig. 12.115. The dynamics of a growth of the shell energy.

In Fig. 12.116, we show the dependence of the surface tension force (in CGS units) on the parameter $\Delta\gamma$ and on the number of a running layer of the target.

The growth of the shell mass leads to the appearance of a braking force. The dependence of the braking force on the number of a layer of the target is shown in Fig. 12.117.

As seen, the total net force acting on the shell accelerates it despite the presence of the braking force.

The motion to the shell center should be described by relativistic equations.

In Figs. 12.118–12.122, we show the dynamics of the motion of a shell in the phase space, namely the dependence of the dimensionless radius of the shell on its running radius. In Fig. 12.118 we present a variant of the dynamics for the almost absent fractality of the shell.

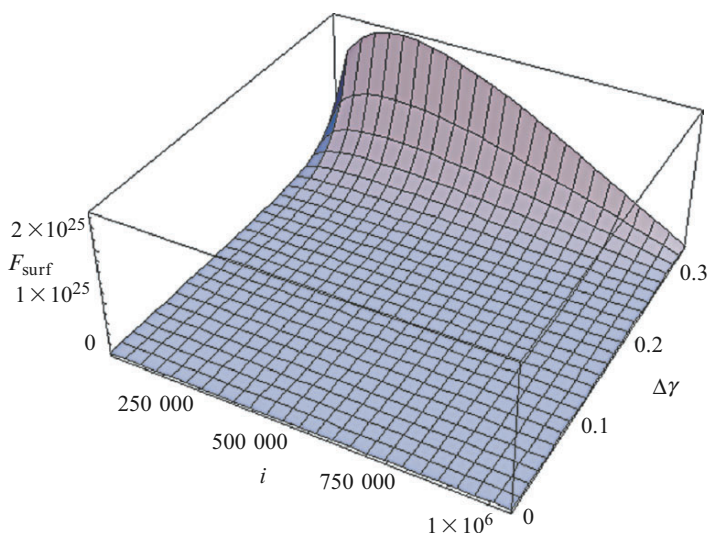


Fig. 12.116. The dependence of the surface tension force (in CGS units) on the parameter $\Delta\gamma$ and on the number of a running layer of the target.

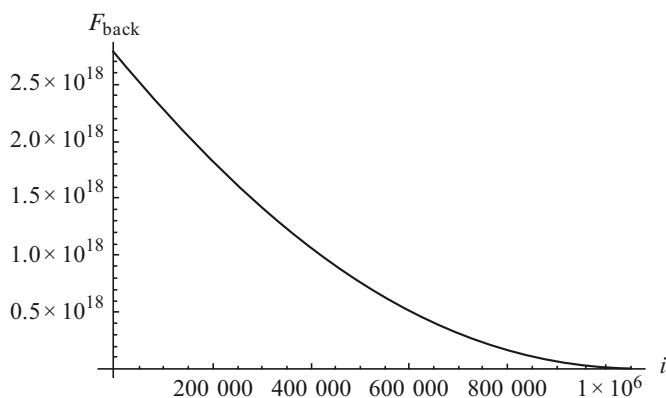


Fig. 12.117. The dependence of the braking force on the number of a layer of the target.

As seen, the shell velocity is nonrelativistic during the whole period of its evolution.

The evolution of the shell momentum defines a change of its kinetic energy while approaching the center.

In Fig. 12.119 we demonstrate the evolution of the kinetic energy of the shell relative to its rest energy.

With increase in the fractality of the shell, the force accelerating it is strongly increased. Already at the parameter of fractality $\Delta\gamma \approx 0.1$, the motion of the shell near the center becomes relativistic (see Fig. 12.120).

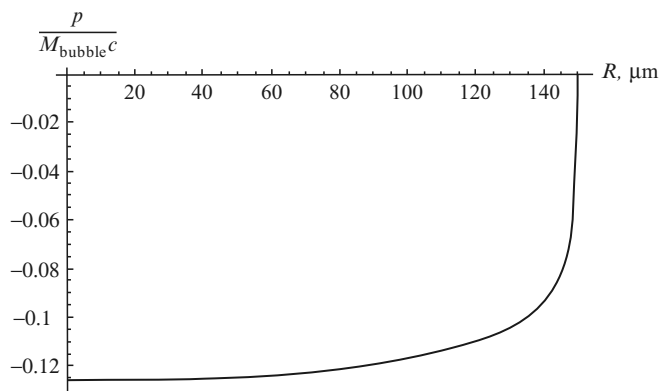


Fig. 12.118. A variant of the dynamics for the almost absent fractality of the shell.

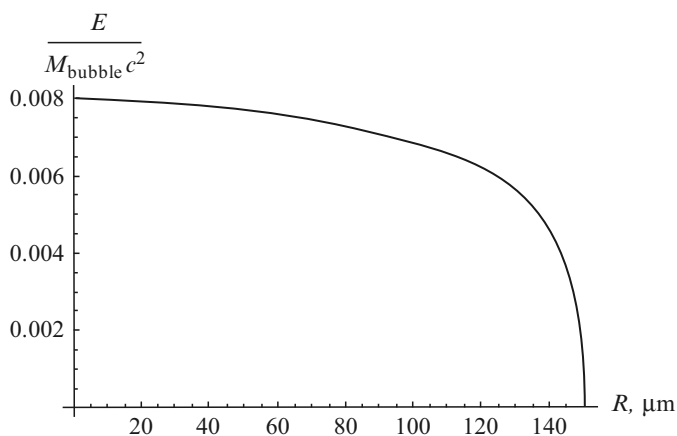


Fig. 12.119. The evolution of the kinetic energy of the shell relative to its rest energy.

The relativistic character of the motion of the shell near the target center corresponds to a sharp increase in its kinetic energy as compared to the case of an almost continuous shell, which is seen in Fig. 12.121.

The structure of the dependence of the force acting on the shell on both the coordinate and the fractality level indicates that the shell moves all the time toward the center. The loss of its stability and the transition to the final stage can occur “in flight”. A reason for the stability loss can consist in a growth of the density and its approach to the nuclear density. In this case, the cluster becomes continuous and loses the stability. The second reason appearing simultaneously with the first one consists in that the relative thickness of the shell increases, and the shell loses the stability due to a

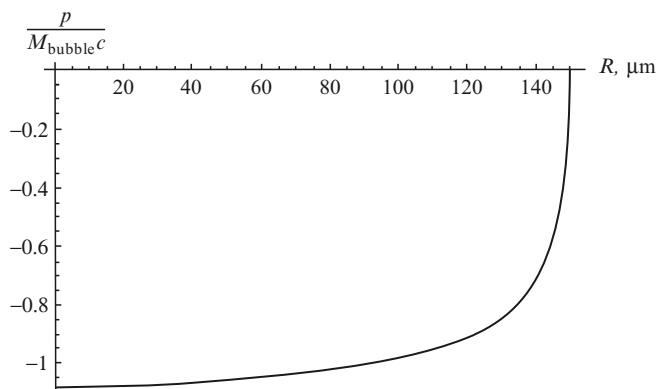


Fig. 12.120. The motion of the shell.

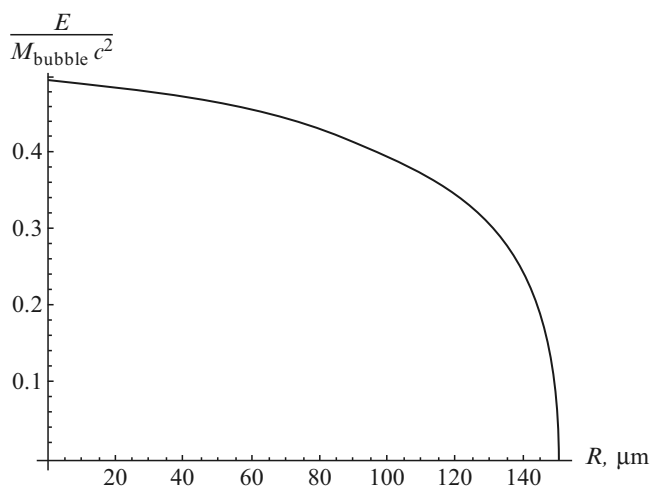


Fig. 12.121. The kinetic energy of the shell.

change of the ratios between the contributions of different components to the binding energy.

At some time moment, all the energy accumulated in the shell E_{bubble} is released, and the shell is dispersed into fragments. In this case, a great kinetic energy is released. Let us consider the energy spent on the formation of fragments, E_{fragment} . This energy can be estimated with the use of the distribution of fragments by mass numbers $f(A)$ after the averaging of the binding energy $B(A)$ of fragments with this distribution:

$$E_{\text{fragment}} = \int_1^{\infty} B(A) f(A) dA.$$

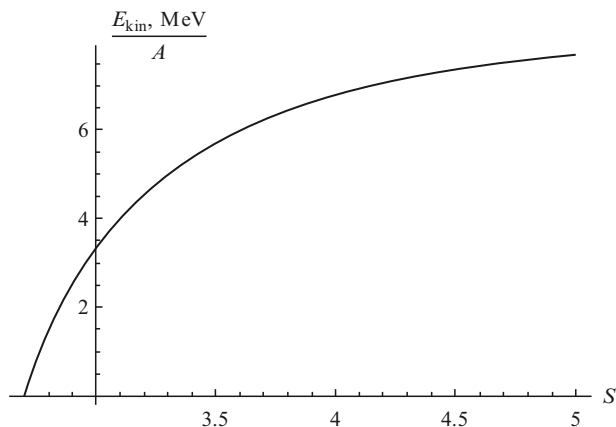


Fig. 12.122. The kinetic energy (in MeV per nucleon) versus the exponent s in the distribution function of fragments by mass numbers $f(A) \propto A^{-s}$.

The kinetic energy E_{kin} is defined by the difference between the energy brought by the shell to the center and the energy necessary for the formation of nuclei dispersed from the target center:

$$E_{\text{kin}} = E(R_b = 0) - E_{\text{fragment}}.$$

In Fig. 12.122, we show the kinetic energy (in MeV per nucleon) versus the exponent s in the distribution function of fragments by mass numbers $f(A) \propto A^{-s}$.

In the example under consideration, the estimation of the parameters of the distribution of fragments by masses gives the value of released kinetic energy to be about 1.2 to 1.6 MeV per nucleon.

12.5. Conclusion

The executed investigation was aimed at the development of basic theoretical ideas which would clarify the experimental results derived at the Electrodynamics Laboratory “Proton-21” and define the further directions for the investigations in the field of nucleosynthesis. One of the attempts to solve the posed problem is described above on the basis of the conception of the interrelation of beam-plasma and nuclear processes conditioned by powerful concentrated flows of energy.

The essence of the above-presented conception consists in the continuity (in the quantitative and qualitative characteristics) of a sequence of physical processes upon the transition from beam-plasma phenomena on macroscales to nuclear phenomena on nuclear scales.

This means that, in order to get a physical interpretation of the processes running in a substance in extreme states, we have designed and realized a scenario of these processes, in which powerful macroscopic actions induce the sequence of phase transitions leading to the appearance of the states of the medium ensuring the possibility to concentrate energy up to great densities on microscales. In this case, the characteristics of the medium approach those of the nuclear substance and allow one to influence the latter.

From the physical viewpoint, these actions lead, upon the mathematical analysis of the whole process, to a change in the initial and boundary conditions for the equations describing the main characteristics of a nucleus. Thus, the results of the action on a nuclear system become dependent on the macroscopic parameters realized in experiments.

Below, we present a brief survey of the results of our investigations and consider the perspectives of the further studies.

As known, the interaction of nuclei is mainly defined by two kinds of forces: namely, the strong short-range interaction (attraction) of nucleons and the long-range electromagnetic interaction (repulsion) of protons (see Ref. 361). Hence, the synthesis of new nuclei from the initial ones or the “repacking” of nucleons is hampered by the Coulomb force of repulsion of protons which belong to different nuclei. That is, to realize the reaction of synthesis of a new nucleus from the interacting nuclei, it is necessary to overcome the Coulomb barrier.

Almost all the attempts to practically execute the reaction of synthesis of nuclei (besides the project realized at the Electrodynamics Laboratory “Proton-21”) were based on the following means to overcome this barrier which seems to be natural: the interacting nuclei should be supplied by the energy of the order of the Coulomb barrier in order to ensure the possibility of their approaching at a distance of the order of 1 fm, which would lead to their further fusion at the expense of the strong interaction. Such approaches to the synthesis of nuclei are based on the viewpoint that the long-range character of Coulomb forces hinders the reaction of synthesis. We note that these approaches are low-efficient, because the formation of a new compound nucleus is inevitably accompanied by its strong excitation, instability, and, finally, decay.

However, just the long-range character of the electromagnetic interaction allows us to propose a completely different method. *The long-range character of the Coulomb forces makes it possible to use collective effects and the processes of self-organization of the nuclear matter*, whose theoretical models were described in our works. These models clarify the sense of the amazing (at first glance) results of experimental studies of the processes

occurring in hard-current diodes performed at the Electrodynamics Laboratory “Proton-21”.

In Chapter 2, the conception and scenario of the self-governing synthesis of nuclei have been presented. Below, we give some details and peculiarities of the scenario by basing on the developed models of collective nuclear processes running in a target under the action of a primary driver [a relativistic electron beam (REB)].

We will estimate the parameters of a diode with copper electrodes which are typical of our experiments. The radius of a target (anode) $R_0 \approx 250 \mu\text{m}$. The epures of the pulses of current and voltage are shown in Fig. 12.29. The maximum difference of potentials is of the order of $U_{\text{max}} = 600 \text{ keV}$, the maximum current is of the order of $I_{\text{max}} = 50 \text{ kA}$, and the pulse duration is about $\Delta t_{\text{beam}} = 20 \text{ ns}$. The energy contained in such a current pulse is of the order of $I_{\text{max}} E_{\text{beam}} = 0.5 U_{\text{max}} I_{\text{max}} \Delta t_{\text{beam}} = 350 \text{ J}$.

As shown in our work, the process of self-governing synthesis of nuclei passes several main stages, and each stage is based on nonlinear collective effects. These stages are the sequence of states basically different one from another, and the transition between these states occurs as a result of the development of the sequence of instabilities.

We note that, due to the initiation of internal self-consistent collective processes, all the stages of the evolution (even those being traditional for the inertial synthesis) involve the appearance of essential features which lead eventually to the self-organization of all basic physical processes occurring in the diode.

The numerical simulation of the processes occurring in the hard-current diode (see Sec. 12.2.1) designed at the Electrodynamics Laboratory “Proton-21” showed that one of its basic peculiarities consists in the dynamics of the cathodic and anodic plasma. *Practically during the entire period of generation of a pulse of electrons, the diode designed at the Electrodynamics Laboratory “Proton-21” is a plasma-based device.*

The numerical and analytic simulation of the processes occurring in the diode showed that, with increase in the difference of potentials applied to the diode, a hard-current beam of electrons is formed due to the explosive emission. Owing to the magnetic self-isolation mode arising in the diode, the beam possesses two components: the component propagating along the system axis and the ring component of the beam.

The interaction of the beam and a target leads to the creation of the anodic plasma covering the target. Practically, the target represents a plasma cone met by the beam (see Fig. 12.27).

The presence of the ring component of a beam leads to that *the region of the interaction of a beam and the lateral surface of the plasma target is*

a thin ring. The calculations showed (see Sec. 12.2.2) that the ring moves rapidly to the cathode on the surface of the plasma cone formed by the anodic plasma. The velocity of the collective motion front exceeds that of individual ions positioned in the same electric field (see Fig. 12.44).

The small angle of inclination of this surface to the cone axis *provides the cumulative effect and increases the pressure on the lateral surface of the target* (see Fig. 12.27).

In this case, the area of the plasma surface of the anode, with which the beam interacts, decreases sharply (at the same time, the area of the plasma cathode increases, on the contrary). Such a dynamics of the emission and transportation of the electron beam results in a sharp (shock-like) concentration of the beam energy on the target surface.

At small currents of the beam, the length $l_\varepsilon(U)$, on which the main share of the electron beam energy is lost in the target, is defined by the instantaneous difference of potentials $U(t)$. This length can be evaluated by Eq. 12.187. For typical values of the voltage maximum, we get $l_\varepsilon(U_{\max}) \approx 300 \mu\text{m}$. However, else prior to the time moment of the attainment of the voltage maximum, the current in the diode turns out to be sufficiently large. For a large current, the typical depth where an electron stops at the expense of collective effects of the dissipation decreases by several times as compared to the depth of penetration of a single electron in the metal. The maximum depth of penetration decreases up to $l_{\varepsilon 0} \leq 20 \mu\text{m}$ by the time moment when the current reaches the critical value I_{cr} (see Ref. 334). Here, I_{cr} is the critical current of electrons for the diode in its near-anodic part and is defined by the constructional features of the diode geometry in the anodic region and by the plasma dynamics in the diode.

Due to the specific structure of the diode, *there occurs the accumulation of a charge in the anode and the locking of the beam*, which depends on the duration of the beam front, maximum current of the beam, and rate of its growth.

The process of locking of a beam begins from the time moment when the beam current reaches the value of the order of I_{cr} . The mechanism of locking is related to a number of nonlinear plasma-beam processes. The effective collision frequency increases with the current as a result of the interaction of the beam with the dense near-surface plasma (see Sec. 12.2.) This results in the deceleration of the beam and then in both the reduction of the length l_ε and *the self-locking of the beam at the expense of the development of the phenomena similar to the formation of a virtual cathode in the diode* (see Ref. 342).

The equations describing these processes are the Poisson equation for the field and the hydrodynamic equations for the flux of electrons

Eqs. 12.192, 12.193. They are solved in Sec. 12.2 numerically and analytically with the above-indicated parameters of the beam. The analysis of the model shows that a reduction of the typical energy absorption length l_ε occurs according to the relation $l_\varepsilon \approx l_{\varepsilon 0} / \sqrt{I(t)/I_{cr}}$, when the beam current approach the critical value I_{cr} . With increase in the supercriticality $I(t)/I_{cr}$, the nonlinearity degree grows, and the length is reduced else rapidly by reaching the micron and submicron values near the current maximum (see Fig. 12.35).

Such an evolution of the absorption length of the beam energy in the near-surface plasma layer causes *the appearance of the self-consistent mode with sharpening, in which the temporal behavior of the beam power absorbed in the anode reveals the explosive character*. This yields that the power attaining the maximum value of the order of $P_{\max} \approx I_{\max} U_{\max} \approx 30$ GW is released near the surface, and almost all the beam energy (~ 300 J) is absorbed in the target volume of the order of $V_{int} \approx l_\varepsilon S_a \leq 10^{-7} \text{ cm}^3$.

This volume contains the number of nucleons $N_A = \rho_A V_{int} / m_p \approx 4 \cdot 10^{17}$ and, respectively, the number of nuclei $N_A = N_{nu} / A \approx 6 \cdot 10^{15}$. The energy per nucleus (ion) is $E_A \approx (300 / N_A) \text{ J}$, i.e., of the order of 320 keV. Because the full ionization of an ion requires the energy of the order of 120 keV, we will deal with the electron-nucleus plasma at the end of the evolution of the beam.

Such a self-consistent self-governing mode with sharpening is essentially different from the standard mode with sharpening, which is used in the inertial nuclear synthesis and involves the explosive type of the evolution of power at the expense of the evolution of the power of a primary driver on the surface of a target. In the standard procedure of synthesis, this circumstance leads to very tough requirements to a primary driver.

In our case due to the realization of the self-consistent mode, the pressure pulse on the surface of the anode (the target) attains very great values (see Sec. 12.2.1) of the order of $p_{\max} \approx 0.5 P_{\max} \Delta t_{beam} / V_{int} \approx (3 \cdot 10^2 / 10^{-7}) (\text{J} / \text{cm}^3)$, i.e., of the order of tens of Gbar (despite the relatively weak energy of a primary driver).

The performed numerical calculations showed that the experiments carried out at the Electrodynamics Laboratory "Proton-21" have demonstrated the record values of pressure exceeding the best modern results (the pressure of the order of several Gbar) in the laser inertial synthesis.

As usual, the concentration of the energy of an external source is characterized by the power surface density. However in our case, it is necessary to create a strongly nonequilibrium situation in the whole thin layer of a substance near the surface. This causes the necessity to introduce the notion of an energy source in the phase space, and the critical quantity is the

power absorbed in a unit volume, where the beam interacts with a target, rather than the power fallen on a unit surface.

We emphasize once more that, *despite the use of electron beams as primary drivers, the processes observed in the cases of the nuclear synthesis with inertial confinement of the plasma and the self-governing nuclear synthesis considered in this work are fundamentally different.* These differences can be seen most clearly if one compares these two scenarios.

The differences appear from the very beginning of the interaction of REB with a target. In the ordinary scheme of the inertial confinement of plasma (see Ref. 217), a nonlinear wave on the surface arises due to the direct action of the power falling on the surface. In our case, even the excitation of a density wave is the strongly nonlinear process which involves a wide spectrum of collective effects in the near-surface layer of the target.

Moreover, in the ordinary scheme, a steepening density wave serves only as a means of both the compression of a substance at the center of the target and the initiation of nuclear processes there. In our case, the mode with sharpening which has arisen in the near-surface layer is a source of both nonlinear density waves in a substance and nonlinear density waves of the volume charge in a dense plasma.

Finally, we indicate one of the essential peculiarities of the process of self-governing nucleosynthesis, namely the appearance of a plasma-field structure of concentric, superdense, superthin shells of a substance which are immersed in the degenerate electron fluid.

A peculiarity of this structure is the strong coherence in the radial direction upon large correlations of states over the surface of this structure.

The analysis of the consequences of the uncertainty relations for the states which are coherent along the radius and reveal the strong correlations on the shell surface was carried out in Sec. 12.3. in the general case. The creation of such coherent and correlated states causes the appearance of the active phase of development of the self-governing synthesis of nuclei. *In this case, the source of energy is, mainly, the mass defect of a cluster upon its fragmentation from the surface.*

We now turn to a more detailed presentation of the main results of the conception constructed by us.

As shown in Sec. 12.3., the quantity characterizing the efficiency of the pulse action on a target is the second time derivative of the specific power, J , absorbed in the surface layer of the target. The experimental data show that the limit value of this quantity, which ensures the successful excitation of nonlinear waves in the target, is $J_{cr} \sim 10^{32} \text{ W}/(\text{cm}^3\text{s}^2)$. The main component of this quantity is the power spent noncoherently (coming into heat) in a unit volume which is $\approx n_A T / \tau_{ee}$ by the order of magnitude.

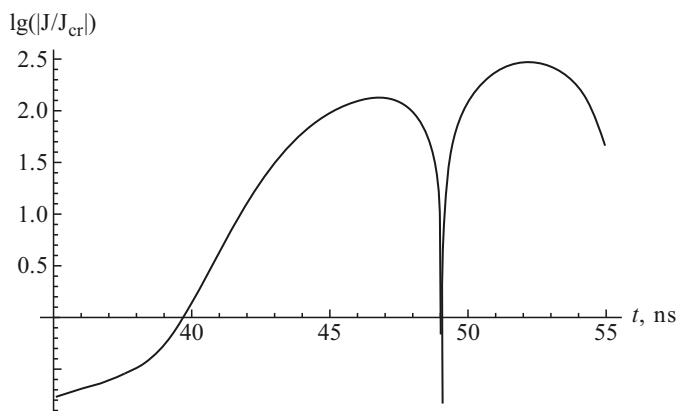


Fig. 12.123. Values of the efficiency criterion for a pulse action versus time.

Just this value of volume density should be exceeded upon the focusing of the beam in order to create a coherent nonequilibrium state. A criterion for the possibility to realize such an action is presented by the ratio J/J_{cr} . The plot of the temporal evolution of the modulus of J/J_{cr} for our driver is given in Fig. 12.123 on the logarithmic scale.

The analysis of Fig. 12.123 indicates that, with the given parameters, the efficiency criterion is satisfied near the region where the current and the energy of the electron beam are maximum, which allows the efficient action on a target to be possible.

The powerful pulse of pressure excites a nonlinear acoustic wave in the self-consistent mode with sharpening. For the power which is supplied by the primary driver to a target, the yield limit of the target substance will be overcome very rapidly, and the medium begins to reveal itself as a viscoelastic one (see Sec. 12.2.2).

On this stage, the evolution of a state of the target is described by Eqs. 12.220–12.224 followed from the one-fluid hydrodynamics with a self-consistent source of pressure of the explosive type on the target surface. These equations reflect the laws of preservation of the energy and momentum in the processes of generation and evolution of the density wave. *With increase in the action power, the substance density wave steepens according to the hydrodynamic equations, which occurs mainly due to the action of hydrodynamic nonlinearities of the type $\vec{u}\nabla\vec{u}$* (see Fig. 12.41, 12.42).

The characteristic distance l_{ex} , which is passed by a wave prior to the steepening (see Ref. 388), can be estimated by a characteristic spatial scale of the high-pressure region on the surface l_ϵ , the isentropic exponent γ , and the Mach acoustic number at the initial time moment upon the excitation

of the nonlinear wave, $M_{ac0} \ll 1$, as

$$l_{ex} \approx R_0 \exp(-1/\sigma_0), \quad \sigma_0 = \frac{\gamma + 1}{2} \frac{2\pi}{l_\varepsilon} R_0 M_{ac0} \quad (12.379)$$

and equals about 100 to 200 μm .

The substance density at the leading edge of the wave is strongly increased (due to the self-consistent mode with sharpening) as compared to the initial density $n_0 \approx 10^{22} \text{ cm}^{-3}$, and the ionization equilibrium shifts rapidly to the region of complete ionization of a substance (see Sec. 12.3.). The substance at the leading edge of the wave becomes the electron-nucleus plasma by the time moment of the steepening.

We note that the conclusion about that arbitrary large values of the substance density are attained, in principle, in the mode with sharpening for a converging spherical or cylindrical wave was drawn in (see Ref. 246). It was hypothesized that the density limitations can arise as a result of the development of an instability. This hypothesis was recently confirmed by the numerical analysis of the processes of laser inertial synthesis in (see Ref. 389). Our calculations showed that the processes with sharpening are indeed limited by an instability. However, this instability (see below) means not the destruction of a closed symmetric acoustic wave which moves to the target center, but the start of a new stage of the evolution. As a result of the development of an instability, the phase transition (to the collapse of Langmuir waves in thin plasma layers) occurs.

The solution of the hydrodynamic equations in the Lagrange variables (see Sec. 12.2.3) showed that the substance density at the leading edge of a wave by the time moment of the development of an instability destroying the mode with sharpening reaches the values of $\approx 10^4 n_0$, and the typical spatial scale of the leading edge $l_f \approx l_{\varepsilon \min}/10^4 \approx 10^{-8} \text{ cm}$ (*in this case, the mean distance between particles in the layer remains to be much less than the thickness of the wave front, and the hydrodynamic equations are not else valid*).

The explosive increase in the substance density on the leading edge of a steepening nonlinear wave of the density is the final stage of the development of the one-fluid hydrodynamic mode. *A further development of the processes in the target is related to the separation of charges and the appearance of a two-fluid mode and a self-consistent electromagnetic field. The two-fluid mode can arise in two ways: as a result of the development of gradient instabilities of the one-fluid mode or due to the accumulation of a volume charge in the thin near-surface layer of the target (from the surface to the region of a virtual cathode in the dense plasma).* These processes are described in Sec. 12.2.4. The evolution of the processes running in the target

by any mechanism leads to the appearance of a nonlinear plasma wave of the volume charge.

From this time moment, the self-consistent electromagnetic fields appear in the condensed medium, and the two-fluid electron-ion hydrodynamic mode of the flow starts.

Like the case of the structures arising directly from the one-fluid mode, the plasma-field structures created upon the development of the nonlinear stage of the two-fluid mode are referred to the type of strongly nonlinear, quasistationary, Langmuir oscillations (see Sec. 12.2.5).

At the estimated density of electrons, the size of a Wigner-Seitz cell is of the order of $R_{WS} \approx 1.5 \text{ to } 2.0 \cdot 10^{-11} \text{ cm}$. Therefore, the layer of nuclei with a typical thickness δ_{str} along the radius contains several such cells.

In fact, every layer includes nuclei which are “swimming” in the degenerate electron fluid with the fluxes of electrons and nuclei penetrating these thin layers-membranes.

The arisen quasistationary plasma-field structure is the system of virtual cathodes and anodes in the electron-nucleus plasma. The virtual electrodes separate the regions with a quasineutral electron-nucleus plasma, which are analogous to domains in a dense solid plasma (see Refs. 346, 359).

In the quasistationary plasma-field structure, certain electromagnetic perturbations are propagating. The kind of an equation describing the spectrum and the dispersion curve correspond to the regularities derived for thin plasma layers and thin beams by virtue of a small thickness of plasma shells.

High-frequency oscillations in the structure of virtual electrodes are similar to the spectra of oscillations in vircathors. We may expect that waves of all kinds propagating in thin layers of plasma interact one with another. Since the electromagnetic interaction has a long-range character, the collective effects change significantly the electromagnetic interaction of charged particles as a result of the excitation of the natural oscillations of the medium (see Sec. 3.3).

The effect of quantization in small-thickness plasma layers in the arisen structure (see Ref. 369 and Sec. 12.3) leads to the appearance of many components with different energies in the plasma (the plasma has a sharp anisotropy of the masses of carriers). The essential difference in the longitudinal and transverse masses of nuclei in a thin plasma layer becomes the reason for the acoustic character of the dispersion of plasmons (the natural plasma oscillations in the plasma-field structure) in the region of small wave numbers:

$$\omega_k = c_{str}k, \quad c_{str} = \omega_{pi}/\chi, \quad \chi = 2\pi/\delta_{str}. \quad (12.380)$$

The interaction of point-like charges representing nuclei is significantly modified in the electron-nucleus plasma at the expense of the contribution of interactions of the wave-particle type. Moreover, the sign of the interaction of nuclei changes, which means the appearance of a phase transition analogous to the phase transition of a substance into the superconducting state. The transition to superconductivity happens as a result of the formation of integer-spin “Cooper pairs” due to the attraction of electrons near the Fermi surface. Just the “Cooper pairs” undergo the Bose–Einstein condensation.

Estimation of the temperature of a phase transition gives a value of $T_c \approx \hbar\omega_p \approx 8$ keV.

A modification of the Coulomb interaction of nuclei as a result of the exchange by plasmons of thin plasma layers and the destruction of the Coulomb barrier between the interacting nuclei cause a sharp increase in the probability of the synthesis of new nuclei.

The last is also affected by the fact that the presence of sources of energy in the “beam–dense plasma of a target” system makes its quasistationary states to be strongly nonequilibrium and endows the energy distributions of particles by power asymptotics (see solutions Eq. 12.296 of the kinetic equations deduced in Sec. 12.3.3). *The increase in the probability of the reaction of fusion of interacting nuclei promotes the growth of a cluster.*

In every separate nuclear layer of each period of the Langmuir nonlinear collapsing structure, being a stationary state arisen as a result of the nonlinear stage of the dynamical evolution of the nuclear matter, a fractal plasma-field structure of the nuclear matter is developed.

The arisen structure has the form of “a nuclear web”. For example, it is shown in Ref. 219 that the base of the structure of heavy nuclei can consist of α -particles composing a nuclear cluster of the fullerene type. The further estimations of the parameters of a cluster were performed for the “web threads” formed by α -particles. These threads have the nuclear scale for diameters and a quite macroscopic scale for their lengths and visible sizes. In our case, the diameters of web threads d_{web} are of the order of several fm, the length l_{web} is about $6 \cdot 10^4$ cm, and the total area of the surface of threads of a cluster is of the order of $S_{\text{web}} \approx 7 \cdot 10^{-8}$ cm². The last value is much more than the surface area $S_{\text{ball}} \approx 10^{-13}$ cm² of a continuous nuclear structure containing the same number of nucleons.

In the general case, the surface of a cluster and the surface of a region, where it is located, are connected by $S/S_0 = 4N^k$, $0 < k < 1/3$, where N is the number of particles in the cluster. The value $k = 0$ corresponds to a continuous cluster, and $k = 1/3$ does to a maximally loose structure. In the

given case, the ratio of the surfaces of a cluster S_{web} and its projection S_0 corresponds to $k \approx 0.252$.

The area of a cluster surface can be expressed directly in terms of the number of particles in the cluster by $S = 4N^\gamma$, where $2/3 \leq \gamma \leq 1$. The quantities k and γ are connected with the fractal dimensionality of a cluster D_f by the relation $k = \gamma - 2/D_f$.

The formation of a macroscopic cluster is accompanied by a modification of the binding energy of nuclei as collective nuclear structures (see Sec. 12.3). These modifications are manifested in all the traditional terms of the Weizsäcker formula. The main changes occur in the surface and Coulomb energies:

- Surface energy $W_s(A, D_f)$ increases in modulus (by remaining negative) as a result of the appearance of a fractal structure of the nuclear cluster surface [$W_s(A, D_f) = -4\pi r_0^2 \sigma_A A^{k+2/D_f}$].
- Negative contribution of the Coulomb volume energy to the binding energy of a cluster decreases significantly (in modulus) due to both the screening of the Coulomb barrier and the effective decrease of the charge of a nuclear cluster, Z_∞ , for the external observer as a result of the condensation of electrons and mesons: $Z_\infty \approx \frac{3}{2} \left(\frac{l}{R}\right)^3 \left(\frac{R}{l} - th \frac{R}{l}\right) A \propto A^{1/3}$.
- Contribution of the symmetry-related energy (the contribution of the weak interaction) to the binding energy of a nucleus varies significantly due to the modification of the surface energy of a cluster.

The appearance of a nuclear cluster in the form of a shell with fractal structure is followed by the active phase of its evolution.

The active phase is related to the fragmentation of a cluster from the surface and the release of energy (to the formation of the mass defect) as a result of the fragmentation and the evolution of the nuclear cluster (see Sec. 12.3.4).

The distribution of fragments $f(A)$ by the number of nucleons which belong to them can be calculated by the theory of “gas-fluid” phase transitions in the nuclear substance and turns out to be a power one (see Sec. 12.3.3):

$$f(A) \propto A^{-\alpha_f}. \quad (12.381)$$

According to the thermodynamic theory (see Sec. 12.3.3) and the Fisher’s drop model (see Ref. 284), we can get the relations for the critical indices. This yields the following estimates: $D_f = \frac{2}{\gamma-k} \geq 2.71$, $2.33 < \alpha_f < 2.5$, and $\sigma_f = 1$. The Fisher’s critical index σ_f defines the dependence of the area of the cluster surface on the number of particles: $S \propto A^{\sigma_f}$. Thus,

at the phase transition under consideration, the area of the cluster surface $\propto A$, which testifies to the fractal character of the structure.

A macroscopic nuclear structure arisen on the certain distances from the surface of a macroscopic nuclear structure moves to the target center at the expense of the energy released upon the evaporation of this structure and is a nonlinear wave with nuclear combustion. We may say (see Sec. 12.3) that the running of a collective two-stage nuclear reaction is accompanied by the formation of a compound nucleus, being a macroscopic fractal nuclear structure. On the first stage of the nuclear reaction, a compound nucleus is formed because the flux of nuclei of the target at the expense of the phase transition leading to the pairing of nuclei enters into the cluster. On the second stage, the compound nucleus (a macroscopic nuclear cluster) evaporates. Since a macroscopic number of nuclei evaporates from its external surface, the reaction yields a very great collection of nuclei, and the process has a statistical character. Due to the collective statistical character of the processes of evaporation, the system evolves according to the variational principles, as it occurs usually in statistical physics. In this case, any time moment is characterized by such a distribution of products of the reaction which ensures the maximum positive gain of energy.

Thus, our theoretical models and experimental results clearly demonstrate that the nuclear processes of synthesis are running slowly (adiabatically) already at the distances to the surface greater than about one hundred microns. Since the reactions of synthesis are running sufficiently slowly, *the probability of the formation of unstable nuclei is very small, and all synthesized nuclei are nonradioactive isotopes of a wide spectrum of elements.*

The surface energy of nuclei, W_s , is proportional to the surface area of a nucleus, S_A , being a drop of the nuclear liquid: $W_s = 4\pi R_A^2 \sigma_A$, $\sigma_A \approx a_a/4\pi r_0 10^{26} \approx 10^{26} \text{ MeV/cm}^2$.

A nuclear structure with large area can accumulate a very great surface energy (of the order of 0.3 MJ).

The energy accumulated as the surface energy of a cluster presents a part of the surface energy of the initial nuclei of a target. We estimated the surface energy of Cu nuclei participating in the reaction of synthesis as $W_s \approx 10 \text{ MJ}$. The volume part of the strong interaction energy of nuclei of the target does not vary practically due to an extremely low compressibility of the substance, which is a consequence of the short-range character of the strong interaction. Moreover, the strong interaction ensures the high strength of “web threads”.

Nuclear structures in the form of thin shells are exceptionally stable against the processes of fission and the β -processes (see Sec. 12.3.3) and work (see Ref. 376), where the stability of nuclei with a completely empty

core was proved up to those with charge Z equal to several hundreds and with the mass number of about 1000). In addition, as shown in Sec. 12.3.4, the binding energy per nucleon for such macroscopic shell nuclear structure with mass number $A > 10^6$ turns out of the order of 1 MeV and does not practically depend on the mass number.

A nuclear cluster, which was formed in the process of self-governing nucleosynthesis and has such a shell structure, is ideally made ready for a controlled expenditure of the stored energy of the substance of nuclei in its composition.

In the active phase, the movement of the shell toward its center is supported, mainly, at the expense of a decrease in the surface energy of the cluster and its fragmentation.

While approaching the center, the system of concentric layers reaches the maximum parameters prior to the loss of stability and the destruction. An important feature of the extreme state of the shell substance is the process of protonization of free neutrons in the shell substance at the attained densities, which is considered in Chapter 1.

Then there occur the dispersion of fragments of the nuclear cluster, which are of various sizes, from the region of the explosion of the wave-shell and the interaction of products of the explosion of the nuclear megacluster with nuclei of the target substance. Upon this interaction, ordinary high-energy nuclear processes are running and can induce the formation of the whole spectrum of products of the reactions (generally including the radioactive products). We recall that, *upon the movement of the wave to the center, the synthesis of nuclei occurs as a result of the adiabatic processes of fragmentation and involves the class of stable nuclei like the phenomenon of cluster radioactivity.*

In Sec. 12.3.4, we have considered the distribution of fragments of a shell, which became unstable, by mass numbers and have compared it with experimental data.

The distribution of fragments by mass numbers was sufficiently thoroughly analyzed for the collisions of heavy ions. In this case, the distribution of fragments by masses has the form $f(A) \approx A^{-2.65}$, and the exponent differs from one presented by the equilibrium thermodynamics of phase transitions. Differences from the thermodynamical theory can be explained with regard to nonextensiveness (see Sec. 12.3.4). The effect of the nonextensiveness of states is significant both for the collisions of heavy nuclei and the decay of a macroscopic nuclear structure.

The distribution of fragments by mass numbers can be determined with the use of the energy distribution over nonextensive states having

the power asymptotics $\propto \exp_q$. Such a generalization yields the following asymptotics of the distribution for great mass numbers:

$$f(A) \propto A^{-\frac{\tau_f}{q-1}}. \quad (12.382)$$

The statistical distribution of nuclei by mass numbers as a result of the fragmentation of a macroscopic nuclear cluster is a consequence of its nonequilibrium state. So, by comparing the thermodynamical estimate and experimental data, the value of the parameter of nonextensiveness $q \approx 1.879$ which agrees with the above-presented estimates of the fractal dimensionality. The imbalance of the system and the existence of correlations between the states of particles of the system lead also to the nonextensiveness of quasistationary distributions of particles by energy.

In Sec. 12.3.3, we have shown that such distributions with power asymptotics at high energies are the solutions of kinetic equations and correspond to the correlations in the system which arise as a result of the existence of the flows of energy or the fluxes of particles. The power asymptotics of the distribution of particles by energy are experimentally revealed in all parts of the spectrum available for measurements. The exponents in the energy distributions are approximately the same in all regions, $f(E) \propto E^{-2.7}$, and correspond to the states with strong correlations studied in Sec. 12.3.

In the experiments performed in our Laboratory (see Refs. 209–212), pulse relativistic electron beams are used as a primary driver.

A characteristic feature of the used driver is that the REB can perturb simultaneously and efficiently the neutral component, electrons, ions, and natural electromagnetic oscillations of the medium. However, the primary driver used now at the Electrodynamics Laboratory “Proton-21” is not uniquely possible.

The main factors inducing the appearance of the plasma-field structure and, as a result, modifying the nuclear interactions are high values of the relevant quantities and their gradients, namely the substance density and the intensity of a self-consistent electric field, being present in this structure. As was mentioned above, these factors are inherent in the models considered in this section and manifest themselves in the interaction of REB and a target. However, it seems to be possible that similar effects and the values of these factors can be derived also with other drivers.

We will list the most interesting, from this viewpoint, drivers which can lead to the development of scenarios similar to the above-described one:

- streamer
- magnetocumulative generator (MCG)

- electromagnetic pulse
- pinch in the unstable dynamical stage

We note that the above-presented main conceptions arisen during the realization of the program aimed at the self-governing nuclear synthesis can be also the basis for solving the problem of controlled thermonuclear synthesis in the traditional schemes. To the maximum degree, this concerns the electrostatic and magnetoelectric confinement of plasma realized in (see Ref. 344) and the electrostatic confinement carried out by the program “Fusor” (see Ref. 343). Both these programs assumed to use plasma-field structures similar to those described in Sec. 12.2.4. We think that the proper choice of both the characteristics of similar setups, as well as their toroidal modifications, and the parameters of the processes running in them will allow one to reach a positive energy yield.

Our position on the way of the formation of new ideas in this region of science, where the ideas of solid-state physics, plasma physics, nuclear physics, etc., are strongly intertwined, is well reflected by the thought presented in Ref. 390.

The above-described preliminary investigations of the self-governing nuclear synthesis, flows in the electron-nucleus plasma, and the created dynamical structures in the nuclear substance and nuclei with regard to their boundedness are only the beginning of the development of a theory of collective processes and self-organization in the nuclear substance. We believe that the first experimental data presented in other sections of this book will become the basis for the development of new ecologically safe nuclear technologies and energy sources.

EPILOGUE

S. V. Adamenko*

The essence of our approach, based on the main hypotheses underlying our conception of the artificial initiation of a plasma wave-shell collapse, consists in the creation of a driver that produces a coherent excitation of the accelerated motion of a set of interacting (differently in the different states) particles by a mass force with the use of a self-focusing electron beam. We call the driver in question a coherent driver, conditionally to a certain degree, because by itself it is not coherent, strictly speaking. But the driver becomes coherent due to the development of the nonlinear processes involving a sharpening of the energy flow on the leading edge of a collapsing wave.

In a first approximation, we succeeded to create such a driver at the Laboratory "Proton-21". In our opinion, just this circumstance allows us confidently to explain the whole totality of "wonders" observed by us, as well as each wonder separately, on the basis of the fundamental axioms of physics, by satisfying the conservation laws and without any "mystical" forces. From our viewpoint, the potentialities of the coherent physical processes should not amaze or puzzle anyone after several decades of the successful development of laser-based technologies. By comparing the impressive laser-induced effects with the observed effects of coherent nuclear processes, we cannot fail to note that the coherent behavior of electrons does not create the measurable mass defects and does not produce free energy, but only redistributes it in the course of time. At the same time, the release of nuclear energy or its absorption are the defining effects in the coherent processes with active participation of nuclei and nucleons forming "transient" dynamical systems. The potential of these effects is, of course, obvious.

It was of course not easy to realize and to formulate the conclusions at which we arrived, with astonishment and doubts, by verifying repeatedly the huge number of observed facts, improbable at first sight. But several thousand experiments and analytic studies executed by a large body of highly skilled professionals with diverse practical experience, as well as a number of testing measurements performed in specialized laboratories of several countries, which uncovered striking effects, compelled us to take responsibility for the conclusions that follow, because we could find no alternative to them:

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1. For the first time, we have brought about the planned and purposeful reproduction of the cosmic self-supporting low-entropy nuclear combustion of substances under laboratory conditions.
2. The possibility to utilize a new nuclear energy source that can satisfy the demands of mankind in the foreseeable future has been demonstrated.
3. The reasonable use of our discovered energy source possibly gives a real chance to prevent the fine natural balances of our planet from being upset by the consumption of ever-increasing amounts of energy.
4. We have proposed a possible tool for the technological neutralization of dangerous wastes, including the radioactive ones, by way of their energetically self-sufficient nuclear combustion.
5. For the first time, we have created and tested a means that allows us to acquire a wide spectrum of long-lived and stable isotopes of transuranium elements; moreover, their amounts are so large that the study of their properties and, possibly, their practical application will become a solvable technical problem.
6. Some of the synthesized long-lived metastable isotopes have nuclear-physical properties that encourage hopes for their use to realize, in the midterm, the production of a high-tech safe nuclear fuel, including fuel for small-scale mobile energy generators.
7. Through the example of artificially initiated collapse in a solid state target, accompanied by the creation of the products of laboratory nucleosynthesis and controlled release of potential energy of the nuclei, the possibilities of unique natural mechanisms of self-organizing complex systems are demonstrated. Thereby, it is shown that these mechanisms inevitably act in response to excitation of the system's free states with sufficient external mass force that initiates the system's evolution in the required direction of its configuration space by a trajectory that corresponds to the principle of dynamic harmonization or that of least action generalized to the case of dynamical systems with modified nonholonomic ties.

The Electrodynamics Laboratory "Proton-21" is open to mutually beneficial collaboration with anyone who is interested in our reports and publications and wants personally or collectively to contribute to their discussion and to the development of the new nuclear physics line of research and technologies.

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