FERMI-LIQUID APPROACH FOR THE DESCRIPTION OF THE INITIAL STAGE OF FRAGMENTATION AT HEAVY NUCLEI COLLISIONS

A.P. IVASHIN, S.V. PELETMINSKII, YU.V. SLYUSARENKO

National Science Center "Kharkiv Institute of Physics and Technology"
(1, Akademichna Str., Kharkiv 61108, Ukraine;
e-mail: ivashin@kipt.kharkov.ua; slusarenko@kipt.kharkov.ua)

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A mechanism is proposed for the initial stage of the instability development that can induce the fragmentation of the nuclear matter arising as a result of collisions of non-relativistic heavy nuclei. The collision of heavy nuclei is simulated as a collision of two unbounded Fermi-liquid "drops". The instability origination in such a system is related to the propagation of increasing oscillations in nuclear matter. These oscillations can exist in a resting Fermi-liquid: the modified Landau zero sound, modified spin and isospin waves, and a combination of these more simple waves. These instabilities are analogous to the beam instability in the ordinary electron plasma. The analysis of the obtained increments of oscillations is performed. Its results can be used for the experimental confirmation of the proposed mechanism of the fragmentation at nuclear collisions. Directions along which nuclear matter "jets" can be expected are specified.

1. Introduction

The fragmentation phenomenon, i.e. the simultaneous decay of the excited nucleus into lighter nuclei (fragments) and separate particles, has been known for a rather long time. However, the scientific interest to this phenomenon has rapidly grown lately due to the experiments on collisions of heavy fast nuclei, which were carried out in different scientific centers of the world (CERN, Switzerland; Dubna, Russia; Berkeley, Oak Ridge, USA; Hamburg, Germany; Orsay, Caen, France). This interest is caused by the prospect to confirm our ideas of the inner structure of nuclei and the character of inter-nucleon interactions and by the possibility to obtain a new fundamental knowledge about the origin of intranuclear and intranucleon forces.

The nuclear matter formed by the collision of heavy nuclei is an unstable object decaying within a very short time. The decay of the arisen nuclear matter is accompanied by the formation of new lighter nuclei and nuclear fragments (the so-called process of nuclear fragmentation). Nowadays, while describing the evolution of the nuclear matter formed after the collision of heavy nuclei, one usually considers two scenarios (see, for example, [1, 2]). According to the first one, the state of statistical equilibrium can be reached in nuclear matter (the thermalization of nuclear matter) permitting the thermodynamic description of the fragmentation process by means of methods of the phase transition theory. According to the second scenario, the resulting nuclear matter is essentially an unstable matter, in which the processes allow only the dynamical description.

In the present paper, we assume the latter scenario. Each colliding nucleus is a system of nucleons, and it can be considered as a generalized Fermi-liquid, whose particles have spin and isospin (distinguishing protons and neutrons) degrees of freedom. Let us note that the Fermi-liquid approach of Landau and Silin [3, 4] is widely applied in theoretical nuclear physics in order to describe the properties of heavy nuclei [5, 6]. In solving the specific problems of theoretical nuclear physics, it is often assumed that heavy nuclei (with $A > 100$) can be considered as infinite drops of the nucleon Fermi liquid in cases where surface effects are unimportant. Then such an assumption enables one to simplify the complicated mathematical treatment without violation of the final results. Note, that a theoretical approach for finite Fermi liquids was elaborated by A.B. Migdal (see [5] and references therein).

Here, we also suppose the nuclear matter, formed due to heavy nuclei collisions, to be a Fermi-liquid consisting of particles with spin and isospin degrees of freedom. The inter-particle interaction in this kind of nuclear liquid is described by such parameters as Landau amplitudes, by analogy with the normal Fermi-liquid theory [3, 4]. So in order to consider the fragmentation formed after the collision, we come to a less complicated problem of the interaction of two drops of the nucleon infinite Fermi-liquid moving relatively to each other. Note that the description of collisions of heavy ions as two drops of infinite quark Fermi-liquids leading to the formation of a quark-gluon plasma was analyzed in [7].

We have demonstrated that the interactions of two nucleon Fermi-liquids in the system can originate a number of instabilities associated with increased oscillations of the modes existing in the ordinary Fermi-liquid. Such waves include a modified Landau zero sound, modified spin waves, waves related to the excitation of isospin degrees of freedom, as well as more complex waves related to a combination of the mentioned simple waves [6].The present instabilities are analogous to the known "beam" instability in the ordinary electron plasma, occurring when charged particle beams pass through the plasma [8, 9]. Since the mentioned instabilities can evolve exactly after the mutual penetration of the colliding Fermi-liquid drops and the formation of a united Fermi-liquid, we consider the formed system as an object which is far from the statistical equilibrium and requires the dynamical description. The dynamics of such a system is described in this work by the kinetic equation for the one-particle distribution function for nucleons in the collisionless approximation which is widely used in the theory of ordinary Fermi-liquid [3, 4]. Such an approach for the phenomena we consider is valid in the case where the characteristic time of instability evolution is less in comparison with the characteristic time of relaxation in our system. In many cases, such an approach allows the analytic description of the initial stage of instability development in the formed nuclear matter without numerical calculations. Note that the use of the mentioned kinetic equation restricts our consideration to non-relativistic nucleon energies. The modification of the kinetic equation for the description of the evolution of the nuclear matter with relativistic nucleons is given in [10, 11]. In these papers, the development of filamentary instabilities in nuclear matter at the stage of the mutual penetration of the two colliding nuclei was considered on the basis of the kinetic equation for relativistic nucleons.

The use of this approach allows us to describe the initial stage of instability development in the nuclear matter formed after the nuclear collision and to obtain the dispersion equations for zero sound oscillations in such a system (the analog of zero sound oscillations in the ordinary Fermi-liquid). The solutions of the equations for small Landau amplitudes are investigated in detail in the case of the collision of nuclei in the normal state, as well as in the case of collisions of heavy excited nuclei. We will find the expressions for the increments of zero-sound oscillations in nuclear matter and show that those increments have characteristic logarithmic dependence on the excitation energy of a nucleus, when the nuclear matter is formed due to the collision of heavy excited nuclei. This characteristic dependence can be used for the experimental proof or disproof of the mechanism of initiation of instabilities in nuclear matter which is suggested in present paper.

We predict the directions (the angular distribution), in which it is possible to expect the matter jets consisting of the fragments of nuclei and separate nucleons and originating as a result of the destruction of nuclei under a collision. In this paper, these propagation directions are related to the direction and value of the relative velocity of two colliding nuclei. The occurrence of such jets is caused, in our opinion, by the further evolution of instabilities, described in this paper. It is known that the experimental exploration of the angular distributions of outgoing matter jets is the important source of data revealing the processes in a formed dense blob of nuclear matter. For this purpose, the special experimental complexes, which allow one to perform the mentioned measurements almost in the whole spatial angle, are created 4π [2,12]. The experimental detection of the mentioned jets in predicted directions would be the evidence for a realization of the suggested mechanism of initiation of instabilities in the colliding nuclei system.

We suppose the confirmation of predictions given in the present paper should be expected, for example, in experiments on nuclear collisions Gd+U or Xe+Sn (INDRA, [2,12]) with incoming nucleus energies of above 145 MeV per nucleon.

2. Basic Kinetic Equations of Two Colliding Fermi-Liquids

A solution of the formulated problem on the kinetics of the collision of two nucleon Fermi-liquids is based on the assumption that the Fermi-liquid energy is a functional $E(f)$ of the one-particle distribution function $f_i(\mathbf{x},\mathbf{p},t)$ of quasiparticles (nucleons; here, the index "i" is used to denote the quasiparticle spin and isospin, $\mathbf{x}$ is a coordinate, $\mathbf{p}$ is the quasiparticle momentum, and t is the time). We want to emphasize that, in the Fermi-liquid theory, the introduced energy functional plays the role analogous to that of a Hamiltonian in the microscopic theory.

The evolution of the one-particle distribution function $f_i(\mathbf{x},\mathbf{p},t)$ in the collisionless approximation ($\omega\tau_r \gg 1$, ω is a frequency, τ_r is a relaxation time) is governed by the kinetic equation

$$\frac{\partial f_i}{\partial t}+\frac{\partial\varepsilon_i}{\partial\mathbf{p}}\frac{\partial f_i}{\partial\mathbf{x}}-\frac{\partial\varepsilon_i}{\partial\mathbf{x}}\frac{\partial f_i}{\partial\mathbf{p}}=0, \qquad (1)$$

where the quasiparticle energy $\varepsilon_i(\mathbf{x},\mathbf{p},f)$ determining the kinetics of the nucleon system represents a variational derivative of the energy functional with respect to the one-particle distribution function

$$\varepsilon_i(\mathbf{x},\mathbf{p},f)=\frac{\delta E(f)}{\delta f_i(\mathbf{x},\mathbf{p})}.$$

A specific form of the quasiparticle energy $\varepsilon_i(\mathbf{x},\mathbf{p},f)$ as a functional of the distribution function in the Fermi-liquid theory is not known. Therefore, let us introduce functions of the quasiparticle interaction F_{ij} (generalized Landau amplitudes) which represent a linear reaction of the quasiparticle energy on small deviations $\delta f_i(\mathbf{x},\mathbf{p},t)$ of the distribution function $f_i(\mathbf{x},\mathbf{p},t)$:

$$\delta\varepsilon_i(\mathbf{x},\mathbf{p},t)=\sum_j\int d\tau^{'}\int d\mathbf{x}^{'}F_{ij}\left(\mathbf{x}-\mathbf{x}^{'};\mathbf{p},\mathbf{p}^{'}\right)\times$$

$$\times\delta f_j\left(\mathbf{x}^{'},\mathbf{p}^{'},t\right), \qquad (2)$$

where $d\tau=d^3p/(2\pi\hbar)^3$. The functions F_{ij} representing the second-order variational derivatives of the energy functional with respect to the one-particle distribution function are main characteristics of the theory that can be experimentally determined (in this connection see, for example, [13, 14]). In view of short-range interaction forces between nucleons, we will consider from here on that the functions F_{ij} have the form

$$F_{ij}\left(\mathbf{x}-\mathbf{x}^{'},\mathbf{p},\mathbf{p}^{'}\right)=F_{ij}\left(\mathbf{p},\mathbf{p}^{'}\right)\delta\left(\mathbf{x}-\mathbf{x}^{'}\right), \qquad (3)$$

where the quantities $F_{ij}\left(\mathbf{p},\mathbf{p}^{'}\right)$ (we will call them also as Landau amplitudes) do not depend on coordinates. Note that assumption (3) is valid only if we neglect the Coulomb forces induced by the proton charge. By linearizing Eq. (1) near a stationary quasi-equilibrium state (see below) described by the distribution function f_{0i}, we can go over to the Fourier components of deviations $\delta f_i=f_i-f_0$ of the one-particle distribution functions from their equilibrium values

$$\widetilde{\delta f}_i(\mathbf{p},\omega,\mathbf{k})=$$

$$=\int d^3x\int\limits_{-\infty}^{\infty}dt\,\delta f_i(\mathbf{x},\mathbf{p},t)\exp(i\omega t-i\mathbf{k}\mathbf{x}). \qquad (4)$$

As a result, using (2), we get the kinetic equation for $\widetilde{\delta f}_i(\mathbf{p},\omega,\mathbf{k})$ as

$$\left(\omega-\mathbf{k}\frac{\partial\varepsilon_{0i}}{\partial\mathbf{p}}\right)\widetilde{\delta f}_i(\mathbf{p},\omega,\mathbf{k})+$$

$$+\mathbf{k}\frac{\partial f_{0i}}{\partial\mathbf{p}}\int d\tau^{'}\sum_j F_{ij}\left(\mathbf{p},\mathbf{p}^{'}\right)\widetilde{\delta f}_i(\mathbf{p}^{'},\omega,\mathbf{k})=0, \qquad (5)$$

where the amplitudes $F_{ij}\left(\mathbf{p},\mathbf{p}^{'}\right)$ are given, in accordance with (2), (3) by the formula

$$F_{ij}\left(\mathbf{p},\mathbf{p}^{'}\right)=\left.\frac{\delta^2E(f)}{\delta f_i(\mathbf{p})\,\delta f_j(\mathbf{p}^{'})}\right|_{f_i=f_{i0}},$$

ε_{0i} is the energy of noninteracting nucleons, and $\varepsilon_0=\frac{p^2}{2m}$ (m is the nucleon mass).

If we assume that the interaction between nucleons is invariant with respect to the spin and isospin transformations, then the structure of Landau amplitudes F for the nucleon liquid will be of the following form:

$$F=I^{(s)}I^{(i)}F^{(0)}+I^{(i)}(\boldsymbol{\sigma\sigma})F^{(s)}+$$

$$+I^{(s)}(\boldsymbol{\tau\tau})F^{(i)}+(\boldsymbol{\sigma\sigma})(\boldsymbol{\tau\tau})F^{(si)}. \qquad (6)$$

Here, $F^{(0)},F^{(s)},F^{(i)},F^{(s,i)}$ are the generalized Landau amplitudes, $I^{(s)}$ and $\boldsymbol{\sigma}$ are, respectively, the unit matrix and Pauli matrices in the spin space, and $I^{(i)}$ and $\boldsymbol{\tau}$ are, respectively, the unit matrix and Pauli matrices in the isospin space. Accordingly in Eq. (5), there will be one of the Landau amplitudes from expression (6). Below, we will consider the oscillations (oscillations of density, spin density, isospin density, or spin-isospin density) depending on these amplitudes. Thereby, hereinafter we will deal with only one linearized kinetic equation

$$\left(\omega - \mathbf{k}\frac{\partial \varepsilon_0}{\partial \mathbf{p}}\right)\widetilde{\delta f}(\mathbf{p},\omega,\mathbf{k})+$$

$$+\mathbf{k}\frac{\partial f_0}{\partial \mathbf{p}}\int d\tau^{'} F\left(\mathbf{p},\mathbf{p}^{'}\right)\widetilde{\delta f}(\mathbf{p},\omega,\mathbf{k}) = 0, \qquad (7)$$

where we will bear in mind $F\left(\mathbf{p},\mathbf{p}^{'}\right)$ as one of the amplitudes from (6).

Now let us obtain the expression for a quasi-equilibrium distribution function from the kinetic equation (7). The equilibrium distribution functions of the resting Fermi-liquid drops will be denoted by $f_0^{(1)}(\mathbf{p})$ and $f_0^{(2)}(\mathbf{p})$. Then in the case where the first drop is in rest, and the second one (impacting) is moving with velocity $\mathbf{u}$, the equilibrium distribution functions will be $f_0^{(1)}(\mathbf{p})$ and $f_0^{(2)}(\mathbf{p}-m\mathbf{u})$, respectively. This implies that the quasi-equilibrium distribution function of two colliding unbounded Fermi-liquids consisting of identical particles has the form

$$f_0(\mathbf{p}) = \alpha_1 f_0^{(1)}(\mathbf{p}) + \alpha_2 f_0^{(2)}(\mathbf{p}-m\mathbf{u}). \qquad (8)$$

Obviously, this function obeys the kinetic equation (1).

We want to explain the meaning of the coefficients α_1 and α_2 appearing in this expression. With this aim, we introduce auxiliary distribution functions $g_0(\mathbf{p})$, $g_0^{(1)}(\mathbf{p})$, and $g_0^{(2)}(\mathbf{p})$ normalized to unity (and, therefore, permitting a probability interpretation) and related to one another analogously to relation (8):

$$g_0(\mathbf{p}) = q_1 g_0^{(1)}(\mathbf{p}) + q_2 g_0^{(2)}(\mathbf{p}-m\mathbf{u}). \qquad (9)$$

From the normalization condition of these functions to unity, it follows that q_1 and q_2 satisfy the equality

$$q_1 + q_2 = 1. \qquad (10)$$

From here we see that the coefficients q_1 and q_2 must be interpreted as the probabilities that an arbitrary selected particle belongs to either a system with the distribution function $g_0^{(1)}(\mathbf{p})$ or to a system with the distribution function $g_0^{(2)}(\mathbf{p})$. However, in this case, the probabilities q_1 and q_2 must be specified by relative frequencies of the events of finding objects with the distribution functions $g_0^{(1)}(\mathbf{p})$ and $g_0^{(2)}(\mathbf{p})$. If we are dealing with two colliding beams of nuclei, then we obtain the following expressions (see (10)) for the probabilities q_1 and q_2:

$$q_1 = N_1/\left(N_1+N_2\right), \quad q_2 = N_2/\left(N_1+N_2\right). \qquad (11)$$

Here, N_1 and N_2 are the densities of nuclei in the first and second beams, respectively.

As was previously noted, the above-given argumentation implies the normalization of the distribution functions to unity. In the case under consideration, this requirement is not fulfilled for the functions $f_0(\mathbf{p})$, $f_0^{(1)}(\mathbf{p})$, and $f_0^{(2)}(\mathbf{p})$ because

$$\int d\tau f_0(\mathbf{p}) = n, \int d\tau f_0^{(1)}(\mathbf{p}) = n_1,$$

$$\int d\tau f_0^{(2)}(\mathbf{p}) = n_2, \qquad (12)$$

where n is the density of the number of particles in the system consisting of two colliding Fermi-liquids; n_1 and n_2 are the densities of the number of particles in each of these Fermi-liquids. But, assuming in (9) that

$$g_0(\mathbf{p}) = \frac{f_0(\mathbf{p})}{n}, \quad g_0^{(1)}(\mathbf{p}) = \frac{f_0^{(1)}(\mathbf{p})}{n_1},$$

$$g_0^{(2)}(\mathbf{p}) = \frac{f_0^{(2)}(\mathbf{p})}{n_2}$$

and introducing the denotations

$$\alpha_1 = nq_1/n_1, \quad \alpha_2 = nq_2/n_2, \qquad (13)$$

we arrive at expression (8).

3. Zero-Sound Dispersion Equation in a Simple Model

In order to solve the kinetic equation (7), we will use a simple model, where the Landau amplitudes do not depend on the momenta $\mathbf{p}$, $\mathbf{p}^{'}$. We also assume that the colliding Fermi-liquids have zero temperature, so that the equilibrium distribution functions of separate drops are given by the formulas

$$f_0^{(1)}(\mathbf{p}) = \theta\left(\varepsilon_{1\mathrm{F}} - \varepsilon\right), f_0^{(2)}(\mathbf{p}) = \theta\left(\varepsilon_{2\mathrm{F}} - \varepsilon\right),$$

$$\varepsilon = \frac{p^2}{2m} = \varepsilon_0(\mathbf{p}), \qquad (14)$$

where $\theta(\varepsilon)$ is the theta function, and $\varepsilon_{1\mathrm{F}}$, $\varepsilon_{2\mathrm{F}}$ are Fermi energies of the colliding drops. It is evident from (14) that, in this case, the difference between the types of nuclei can be found only in their different Fermi energies (note that, for heavy nuclei, $\varepsilon_{1\mathrm{F}} \approx \varepsilon_{2\mathrm{F}}$).

Within the used assumptions, the linearized kinetic equation (7) can be written as

$$\widetilde{\delta f}(\mathbf{p},\omega,\mathbf{k})(\omega-\mathbf{kv})+$$

$$+\mathbf{k}\frac{\partial f_0(\mathbf{p})}{\partial \mathbf{p}}F\int d\tau \widetilde{\delta f}(\mathbf{p},\omega,\mathbf{k})=0, \quad (15)$$

where $\mathbf{v}=\dfrac{\partial\varepsilon}{\partial\mathbf{p}}=\dfrac{\mathbf{p}}{m}$, and, in accordance with (8) and (14), we get

$$\frac{\partial f_0}{\partial\mathbf{p}}=-\alpha_1\mathbf{v}\delta(\varepsilon_{1\mathrm{F}}-\varepsilon)-\alpha_2(\mathbf{v}-\mathbf{u})\times$$

$$\times\delta\left(\varepsilon_{2\mathrm{F}}-\frac{m(\mathbf{v}-\mathbf{u})^2}{2}\right). \quad (16)$$

The most general solution of Eq. (15) looks as

$$\widetilde{\delta f}(\mathbf{p},\mathbf{k},\omega)=-F\{\omega-\mathbf{kv}+io\}^{-1}\mathbf{k}\frac{\partial f_0(\mathbf{p})}{\partial\mathbf{p}}\delta\varphi(\omega,\mathbf{k})+$$

$$+\delta A(\mathbf{p},\mathbf{k})\delta(\omega-\mathbf{vk}), \quad (17)$$

where we denoted

$$\delta\varphi(\omega,\mathbf{k})\equiv\int d\tau\widetilde{\delta f}(\mathbf{p},\mathbf{k},\omega), \quad (18)$$

and $\delta A(\mathbf{p},\mathbf{k})$ are arbitrary functions satisfying the condition (see (4))

$$\delta A^*(\mathbf{p},\mathbf{k})=\delta A(\mathbf{p},-\mathbf{k}). \quad (19)$$

Formula (17) allows us to find a value of $\delta\varphi(\omega,\mathbf{k})$ in terms of the functions $\delta A(\mathbf{p},\mathbf{k})$:

$$\delta\varphi(\omega,\mathbf{k})=\tilde\varepsilon^{-1}(\omega,\mathbf{k})\widetilde{\delta A}(\omega,\mathbf{k}), \quad (20)$$

where (see(19))

$$\widetilde{\delta A}(\omega,\mathbf{k})=\int d\tau\delta A(\mathbf{p},\mathbf{k})\delta(\omega-\mathbf{kv}), \quad (21)$$

$$\widetilde{\delta A}^*(\omega,\mathbf{k})=\widetilde{\delta A}(-\omega,-\mathbf{k}).$$

Substituting further (20) into (17), we have the relation

$$\widetilde{\delta f}(\mathbf{p},\mathbf{k},\omega)=$$

$$=\delta A(\mathbf{p},\mathbf{k})\delta(\omega-\mathbf{kv})-F\{\omega-\mathbf{kv}+io\}^{-1}\times$$

$$\times\tilde\varepsilon^{-1}(\omega,\mathbf{k})\widetilde{\delta A}(\omega,\mathbf{k})\mathbf{k}\frac{\partial f_0(\mathbf{p})}{\partial\mathbf{p}}, \quad (22)$$

where $\tilde\varepsilon(\omega,\mathbf{k})$ in (20), (22) is defined by the formula

$$\tilde\varepsilon(\omega,\mathbf{k})=\tilde\varepsilon^*(-\omega,-\mathbf{k})=\tilde\varepsilon_1(\omega,\mathbf{k})+i\tilde\varepsilon_2(\omega,\mathbf{k})=$$

$$=1+F\mathbf{k}\int d\tau\frac{\partial f_0(\mathbf{p})}{\partial\mathbf{p}}\{\omega-\mathbf{kv}+io\}^{-1}. \quad (23)$$

In the case of a charged Fermi-liquid, $\tilde\varepsilon$ represents a complex dielectric permittivity (see, for example, [15]). As is known, the presence of an imaginary additional component in $\tilde\varepsilon(\omega,\mathbf{k})$ indicates the damping or increase of wave amplitudes. The wave dispersion law $\omega(k)=\omega_0(\mathbf{k})+i\gamma(\mathbf{k})$ should be found from the equation

$$\tilde\varepsilon(\omega_0(\mathbf{k})+i\gamma(\mathbf{k}),\mathbf{k})=0 \quad (24)$$

with the decrement (increment) $\gamma(\mathbf{k})$ determined by the imaginary part of $\tilde\varepsilon(\omega,\mathbf{k})$. This is the reason why weakly damping or weakly increasing oscillations in the system with

$$|\omega_0(\mathbf{k})|\gg|\gamma(\mathbf{k})| \quad (25)$$

can exist only under the condition that

$$|\tilde\varepsilon_1(\omega,\mathbf{k})|\gg|\tilde\varepsilon_2(\omega,\mathbf{k})|. \quad (26)$$

Using the formula

$$(z+io)^{-1}=P\frac{1}{z}-i\pi\delta(z)$$

(where P stands for the principal value) and taking into account (3.), we write functions $\tilde\varepsilon_1(\omega,k)$ and $\tilde\varepsilon_2(\omega,k)$ (see (23)) in the following form:

$$\tilde\varepsilon_1(\omega,k)=1+\alpha_1\mathcal{F}\left\{1-\frac{\omega}{2kv_{1\mathrm{F}}}\ln\left|\frac{\omega+kv_{1\mathrm{F}}}{\omega-kv_{1\mathrm{F}}}\right|\right\}+$$

$$+\frac{\alpha_2 v_{2\mathrm{F}}\mathcal{F}}{v_{1\mathrm{F}}}\left\{1-\frac{\omega-ku\cos\alpha}{2kv_{2\mathrm{F}}}\ln\left|\frac{\omega-ku\cos\alpha+kv_{2\mathrm{F}}}{\omega-ku\cos\alpha-kv_{2\mathrm{F}}}\right|\right\},$$

$$\tilde\varepsilon_2(\omega,k)=\frac{\pi}{2}\mathcal{F}\left\{\alpha_1\frac{\omega}{kv_{1\mathrm{F}}}\theta\left(1-\left|\frac{\omega}{kv_{1\mathrm{F}}}\right|\right)+\right.$$

$$+\alpha_2\left(\frac{v_{2\mathrm{F}}}{v_{1\mathrm{F}}}\right)\frac{\omega-ku\cos\alpha}{kv_{2\mathrm{F}}}\theta\left(1-\left|\frac{\omega-ku\cos\alpha}{kv_{2\mathrm{F}}}\right|\right)\bigg\}. \quad (27)$$

Here, $\mathcal{F}$ is the dimensionless Landau amplitude,

$$\mathcal{F}\equiv\frac{Fm^2v_{1\mathrm{F}}}{2\pi^2\hbar^3}, \quad (28)$$

velocities $v_{1\mathrm{F}}, v_{2\mathrm{F}}$ are given by the relations

$$\varepsilon_{1\mathrm{F}}=\frac{mv_{1\mathrm{F}}^2}{2},\quad \varepsilon_{2\mathrm{F}}=\frac{mv_{2\mathrm{F}}^2}{2} \quad (29)$$

and the angle α is an angle between the directions of the wave vector $\mathbf{k}$ and the impacting drop velocity $\mathbf{u}$. In accordance with (24) – (27), a dispersion equation has the form $\tilde{\varepsilon}_1(\omega_0(k),k)=0$ or

$$1+\alpha_1\mathcal{F}\left\{1-\frac{s}{2}\ln\left|\frac{s+1}{s-1}\right|\right\}+\alpha_2\frac{1}{\eta}\mathcal{F}\times$$
$$\times\left\{1-\frac{1}{2}(\eta s-s_0)\ln\left|\frac{\eta s-s_0+1}{\eta s-s_0-1}\right|\right\}=0, \quad (30)$$

and the decrement (increment) is given by

$$\gamma(k)=\left\{\frac{\partial\tilde{\varepsilon}_1(\omega,k)}{\partial\omega}\right\}^{-1}_{\omega=\omega_0}\tilde{\varepsilon}_2(\omega,k), \quad (31)$$

$$\tilde{\varepsilon}_2(\omega,k)=\frac{\pi}{2}\mathcal{F}\left\{\alpha_1 s\theta(1-|s|)+\right.$$
$$\left.+\alpha_2\frac{1}{\eta}(\eta s-s_0)\theta(1-|\eta s-s_0|)\right\}.$$

We use the following denotations in formulas (30), (31) (see also (28), (29)):

$$s\equiv\frac{\omega_0}{kv_{1\mathrm{F}}},\quad \eta\equiv\frac{v_{1\mathrm{F}}}{v_{2\mathrm{F}}},\quad s_0\equiv\frac{u}{v_{2\mathrm{F}}}\cos\alpha \quad (32)$$

(remember that, for the heavy nuclei, $\varepsilon_{1\mathrm{F}}\approx\varepsilon_{2\mathrm{F}}$, that is, $\eta\approx 1$).

The derivative $\left\{\frac{\partial\tilde{\varepsilon}_1(\omega,k)}{\partial\omega}\right\}_{\omega=\omega_0}$, in (31), according to (27), (30), acquires the form

$$\left\{\frac{\partial\tilde{\varepsilon}_1(\omega,k)}{\partial\omega}\right\}_{\omega=\omega_0}=\frac{s}{\omega_0}\left\{-\frac{\alpha_1\mathcal{F}}{2}\ln\left|\frac{s+1}{s-1}\right|+\frac{\alpha_1\mathcal{F}s}{s^2-1}-\right.$$
$$\left.-\frac{\alpha_2\mathcal{F}}{2}\ln\left|\frac{\eta s-s_0+1}{\eta s-s_0-1}\right|+\alpha_2\mathcal{F}\frac{\eta s-s_0}{(\eta s-s_0)^2-1}\right\}. \quad (33)$$

4. Solution of the Dispersion Equation for Zero-Sound Oscillations

A dispersion law of the zero-sound oscillations $\omega=\omega_0(k)$ is determined by the solution of Eq. (30). Note first of all that this equation (in the same way as the relation for $\gamma(k)$, see (31), (33)) is invariant with respect to the simultaneous substitution $s\to -s$, $s_0\to -s_0$, (see (32)). This condition shows that the propagation of a zero-sound wave is possible in opposite directions with the same increment or decrement. Consequently, let us consider $s>0$ for the sake of definiteness. Since Eq. (30) is quite complicated in its general form, we will make certain assumptions in order to obtain a solution.

It is known that, in the ordinary Fermi-liquid, the propagation of undamped zero-sound oscillations at zero temperature is possible only under the condition $\mathcal{F}>0$, $s>1$ (in this connection, see [16, 17]). Otherwise, oscillations will quickly damp. Let us make an assumption that

$$\mathcal{F}>0,\quad s>1. \quad (34)$$

We will demonstrate that, in this case, oscillations can increase or damp due to the existence of the impacting drop. Naturally, one can easily see from (31) that the possibility for weakly increased or weakly damped oscillations to propagate in the case determined by inequalities (34) is related to the fulfilment of the condition

$$|\eta s-s_0|<1. \quad (35)$$

If inequalities (34), (35) are valid, the quantity $\tilde{\varepsilon}_2(\omega_0,k)\equiv\tilde{\varepsilon}_2(s)$ has the form

$$\tilde{\varepsilon}_2(s)=\alpha_2\frac{\pi}{2\eta}(\eta s-s_0)\mathcal{F}. \quad (36)$$

Let us study the solution of the dispersion equation (30) with small Landau amplitudes, $\mathcal{F}$, $\mathcal{F}\ll 1$, and under conditions (34), (35). Since, in this case, the quantity s is close to unity, $s\approx 1+\delta s$, Eq. (30) can be written as

$$1+\alpha_1\mathcal{F}+\frac{\alpha_2}{\eta}\mathcal{F}-\frac{\alpha_2}{2\eta}\mathcal{F}(\eta-s_0)\times$$
$$\times\ln\left\{\frac{1+(\eta-s_0)}{1-(\eta-s_0)}\right\}=\frac{\alpha_1}{2}\mathcal{F}\ln\frac{2}{\delta s},$$

whereof

$$\delta s=2\left\{\frac{1+(\eta-s_0)}{1-(\eta-s_0)}\right\}^{\frac{\alpha_2}{\alpha_1}(\eta-s_0)}\times$$
$$\times\exp\left\{-2\left(1+\frac{\alpha_2}{\alpha_1\eta}\right)\right\}\exp\left\{-\frac{2}{\alpha_1\mathcal{F}}\right\}, \quad (37)$$

where

$s = 1 + \delta s.$

Taking into account (33), we have, under conditions (34), (35), the relations

$$\left\{\frac{\partial \tilde{\varepsilon}_1}{\partial \omega}\right\}_{\omega=\omega_0} \approx \frac{1}{\omega_0}\alpha_1 \frac{\mathcal{F}}{2\delta s}\ , \quad \omega_0 = kv_{1\mathrm{F}}\ .$$

Further, by using formulas (31), (36), and (37), we get the following relation for $\gamma(k)$:

$$\frac{\gamma(k)}{\omega_0} \approx 2\pi \frac{\alpha_2}{\eta\alpha_1}(\eta - s_0) \times \times \left\{\frac{1+(\eta-s_0)}{1-(\eta-s_0)}\right\}^{\frac{\alpha_2}{\alpha_1}(\eta-s_0)} \times \times \exp\left\{-\frac{2}{\alpha_1\mathcal{F}}\left(1+\alpha_1\mathcal{F}+\frac{\alpha_2\mathcal{F}}{\eta}\right)\right\}, \quad (38)$$

$$|\eta - s_0| < 1, \quad \omega_0 = kv_{1\mathrm{F}}, \quad \mathcal{F} \ll 1.$$

According to (4) and (24), the condition $\gamma(k) > 0$ or $0 < (\eta - s_0) < 1$ corresponds to the damping of oscillations, whereas the condition $\gamma(k) < 0$ or $-1 < (\eta - s_0) < 0$ corresponds to the increase of oscillations, i.e., the instability development. The latter inequality can be written (in view of (32)) as $\eta < s_0 < \eta + 1$ or, taking into account that $\eta = \sqrt{\varepsilon_{1\mathrm{F}}/\varepsilon_{2\mathrm{F}}}$, as

$$v_{1\mathrm{F}} < u\cos\alpha < v_{1\mathrm{F}} + v_{2\mathrm{F}}, \quad (39)$$

where α is an angle between the vectors $\mathbf{u}$ and $\mathbf{k}$.

It is evident from (38) that, when approaching the right boundary of interval (39), i.e., if

$$s_0 \to \eta + 1\ , \quad s = 1 + \delta s, \quad (40)$$

the increment $\gamma(k)$ is unrestrictedly increasing. However, it is evident that, in case (40), we need to solve the more correct dispersion equation (30). In this approximation and keeping $s \to 1$ in mind, it takes the form

$$1 + \alpha_1\mathcal{F} + \frac{\alpha_2}{\eta}\mathcal{F} = -\frac{1}{2}\alpha_1\mathcal{F}\ln\frac{\delta s}{2} - \frac{\alpha_2}{2\eta}\mathcal{F}\ln\frac{\eta\delta s}{2}.$$

The solution of this equation is given by the relation

$$\delta s = 2\eta^{-\dfrac{\alpha_2}{\eta_1\alpha_1+\alpha_2}} \times \times \exp\left\{\frac{-2\eta}{(\eta\alpha_1+\alpha_2)\mathcal{F}}\left[1+\alpha_1\mathcal{F}+\frac{\alpha_2\mathcal{F}}{\eta}\right]\right\}. \quad (41)$$

We indicate that, according to (33) under conditions (31), (36), (40), the relation for $\left\{\frac{\partial \tilde{\varepsilon}_1}{\partial \omega}\right\}_{\omega=\omega_0}$ is

$$\left\{\frac{\partial \tilde{\varepsilon}_1}{\partial \omega}\right\}_{\omega=\omega_0} \approx \frac{1}{\omega_0}\frac{\eta\alpha_1+\alpha_2}{\eta}\frac{\mathcal{F}}{2\delta s}.$$

Now, in view of (31), (36), and (4.), we can obtain the following formula for the increment $\gamma(k)$:

$$\frac{\gamma(k)}{\omega_0} \approx -2\pi\frac{\alpha_2}{\eta\alpha_1+\alpha_2}\eta^{-\dfrac{\alpha_2}{\eta\alpha_1+\alpha_2}} \times \times \exp\left\{\frac{-2\eta}{(\eta\alpha_1+\alpha_2)\mathcal{F}}\left[1+\alpha_1\mathcal{F}+\frac{\alpha_2\mathcal{F}}{\eta}\right]\right\}, \quad (42)$$

$$\mathcal{F} \ll 1, \quad \omega_0 \approx kv_{1\mathrm{F}}, \quad s \approx 1, \quad s_0 \lesssim \eta + 1.$$

Here, the quantities s_0, η, s are defined as before by relations (32). Taking into account that $(\eta\alpha_1/\eta\alpha_1 + \alpha_2) < 1$ and comparing relations (38) and (42), one can easily conclude that, when $\mathcal{F} \ll 1$, the maximum value of the increment is reached for the upper limit of interval (39), i.e., when

$$\sqrt{E}\cos\alpha \lesssim \sqrt{\varepsilon_{1\mathrm{F}}} + \sqrt{\varepsilon_{2\mathrm{F}}}, \quad (43)$$

where E is the kinetic energy per one nucleon in the impacting drop.

5. Temperature Impact

So far we have discussed the problem of the propagation of weakly increasing or damping zero-sound oscillations, assuming that the colliding Fermi-liquid drops have zero temperature (see (14)). However, when two equilibrium Fermi-liquid drops collide and the temperature of each one is non-zero, a scenario of the weak instability development can have significant peculiarities.

In order to solve this problem, we need to modify relations (30) and (31) determining the frequency and increments (decrements) of zero-sound oscillations in conformity with the case of the colliding equilibrium Fermi-liquid drops with the non-zero temperature. In this case, the equilibrium distribution functions of the drops are (see (8)—(13))

$$f_0^{(1)}(\mathbf{p}) = \left\{\exp\left[\left(\frac{\mathbf{p}^2}{2m} - \varepsilon_{1\mathrm{F}}\right)\Big/T_1\right] + 1\right\}^{-1},$$

$$f_0^{(2)}(\mathbf{p}-m\mathbf{u})=$$
$$=\left\{\exp\left[\left(\frac{(\mathbf{p}-m\mathbf{u})^2}{2m}-\varepsilon_{2\mathrm{F}}\right)\Big/T_2\right]+1\right\}^{-1}, \qquad (44)$$

where T_1 and T_2 are the temperatures of the resting and impacting drops, respectively. For the drops, the relations

$$(T_1/\varepsilon_{1\mathrm{F}})\ll 1, \quad (T_2/\varepsilon_{2\mathrm{F}})\ll 1 \qquad (45)$$

should hold true. These relations determine the possibility to apply the Fermi-liquid description to systems of strong interacting fermions (in the present case - nucleons). Taking into account formulas (44) and inequalities (45), one can make the low temperature expansion in expression (23), which results in the following equation for the determination of the frequency $\omega_0 = skv_{1\mathrm{F}}$ of zero-sound oscillations in the main approximation in the parameters $(T_1/\varepsilon_{1\mathrm{F}})$, $(T_2/\varepsilon_{2\mathrm{F}})$ (see [18]):

$$\varepsilon_1(s)\equiv\tilde{\varepsilon}_{10}(s)+\tilde{\varepsilon}_{1T}(s)=0,$$

$$\tilde{\varepsilon}_{10}(s)=1+\alpha_1\mathcal{F}\left\{1-\frac{s}{2}\ln\left|\frac{s+1}{s-1}\right|\right\}+$$

$$+\alpha_2\frac{1}{\eta}\mathcal{F}\left\{1-\frac{1}{2}(\eta s-s_0)\ln\left|\frac{\eta s-s_0+1}{\eta s-s_0-1}\right|\right\}=0,$$

$$\tilde{\varepsilon}_{1T}(s)=\frac{-\alpha_1\mathcal{F}}{2}P\int\limits_0^\infty\frac{zdz}{e^z+1}\frac{2+(s^2-1)}{(s^2-1)^2(\varepsilon_{1\mathrm{F}}/T_1)^2-z^2}-$$

$$-\frac{\alpha_2\mathcal{F}}{2\eta}P\int\limits_0^\infty\frac{zdz}{e^z+1}\frac{2+\left[(\eta s-s_0)^2-1\right]}{\left[(\eta s-s_0)^2-1\right](\varepsilon_{2\mathrm{F}}/T_2)^2-z^2}. \qquad (46)$$

Here, P means, as before, the symbol of the principal value. In the main approximation, the expression for the quantity $\tilde{\varepsilon}_2(s)$ defining, in accordance with (27), (31), a weak increase or damping of zero-sound oscillations can be put as

$$\tilde{\varepsilon}_2(s)=\frac{\pi}{2}\mathcal{F}\left\{\alpha_1 s f_0\left(\frac{\varepsilon_{1\mathrm{F}}}{T_1}(s^2-1)\right)+\right.$$

$$\left.+\frac{\alpha_2(\eta s-s_0)}{\eta}f_0\left(\frac{\varepsilon_{2\mathrm{F}}}{T_2}\left[(\eta s-s_0)^2-1\right]\right)\right\}, \qquad (47)$$

where

$$f_0(\varepsilon/T)=\left\{e^{\varepsilon/T}+1\right\}^{-1}. \qquad (48)$$

It is easy to see that, at $T_1\to 0$, $T_2\to 0$, the quantity $\tilde{\varepsilon}_{1T}(s)$ turns to zero. Hence, Eq. (46) transforms into the dispersion equation (30). Expression (47), with account of (48), transforms to formula (31) for the quantity $\tilde{\varepsilon}_2(s)$ determining the instabilities of oscillations, when the drops have zero temperature. As we can see from (46), (47), the temperature effects can greatly impact the development of the instabilities associated with the propagation of zero-sound oscillations in the system, when $s\approx 1$, i.e., $\mathcal{F}\ll 1$. Keeping this fact in mind, let us analyze the possibility of a weak increase of zero-sound oscillations in the system, when the resting Fermi-liquid drop is found at zero temperature, T_1=0. Here, we also assume that the following relations (compare with (34), (35), (40), (43), (45)) are valid:

$$\mathcal{F}\ll 1, \quad s\gtrsim 1, \quad s_0\lesssim \eta+1, \quad (\varepsilon_{2\mathrm{F}}/T_2)\gg 1,$$

$$\delta s\,(\varepsilon_{2\mathrm{F}}/T_2)\ll 1, \quad \delta s=s-1\ll 1. \qquad (49)$$

Under these conditions, the dispersion equation (46) meeting these conditions in the main approximation in δs will take the form

$$\tilde{\varepsilon}_1(s)\approx 1+\alpha_1\mathcal{F}+\alpha_2\mathcal{F}(1-I_0)-$$

$$-\frac{\alpha_2\mathcal{F}}{\eta}\ln\frac{4\varepsilon_{2\mathrm{F}}}{T_2}+\frac{\alpha_1\mathcal{F}}{2}\ln\frac{\delta s}{2}=0. \qquad (50)$$

In order to obtain this equation, we used the asymptotic estimate of the integral, when $|t|\ll 1$:

$$P\int\limits_0^\infty\frac{zdz}{(e^z+1)(t^2-z^2)}\underset{|t|\ll 1}{\longrightarrow}I_0+\frac{1}{2}\ln|t|-I_1t^2, \qquad (51)$$

$$I_0=\frac{1}{4}-\int\limits_1^\infty\frac{dz}{z(e^z+1)}+\frac{1}{2}\int\limits_0^1\frac{dz}{z}\left\{\tanh\frac{z}{2}-\frac{z}{2}\right\}\approx 0.07,$$

$$I_1=\int\limits_1^\infty\frac{dz}{z^3(e^z+1)}-\frac{1}{2}\int\limits_0^1\frac{dz}{z^3}\left\{\tanh\frac{z}{2}-\frac{z}{2}\right\}\approx 0.11.$$

The integrals having exactly the same form describe the temperature impact on the dispersion equation (46). The solution of Eq. (50) looks as

$$\delta s = 2\left(4\frac{\varepsilon_{2\mathrm{F}}}{T_2}\right)^{\frac{\alpha_2}{\eta\alpha_1}} \times \exp\left\{-\frac{2Q_1}{\alpha_1\mathcal{F}}\right\}, \tag{52}$$

$$Q_1 \equiv 1 + \alpha_1\mathcal{F} + \frac{\alpha_2\mathcal{F}}{\eta}(1 - I_0).$$

In accordance with (49), this solution must satisfy the relation $\delta s\,(\varepsilon_{2\mathrm{F}}/T_2) \ll 1$. Whence, remembering that $(\varepsilon_{2\mathrm{F}}/T_2) \gg 1$, it is easy to determine a "temperature" range, when solution (52) exists for δs:

$$1 \ll \frac{\varepsilon_{2\mathrm{F}}}{T_2} \ll \exp\left\{\frac{2\eta Q_1}{(\eta\alpha_1 + \alpha_2)\mathcal{F}}\right\}. \tag{53}$$

Noting further that, under (50), the formula

$$\left\{\frac{\partial\tilde{\varepsilon}_1}{\partial\omega}\right\}_{\omega=\omega_0} \approx \frac{1}{\omega_0}\frac{\alpha_1\mathcal{F}}{2\delta s}, \quad \omega_0 = kv_{1\mathrm{F}}$$

holds and using the fact that, when $T_1 \to 0$ and $s \gtrsim 1$,

$$\tilde{\varepsilon}_2 = -\frac{\pi}{4}\frac{\alpha_2}{\eta}\mathcal{F}, \tag{54}$$

we arrive at the following expression for the increment $\gamma_k(s)$, considering (31):

$$\frac{\gamma_k}{\omega_0} = -\pi\frac{\alpha_2}{\eta\alpha_1}\left(4\frac{\varepsilon_{2\mathrm{F}}}{T_2}\right)^{\frac{\alpha_2}{\eta\alpha_1}} \times$$
$$\times \exp\left\{-\frac{2}{\alpha_1\mathcal{F}}\left[1 + \alpha_1\mathcal{F} + \frac{\alpha_2\mathcal{F}}{\eta}(1 - I_0)\right]\right\}. \tag{55}$$

It is true together with conditions (49), (53). In order to obtain formula (54), we have used expression (47) and also have taken into account that (see (48))

$$f_0\left(\frac{\varepsilon_{1\mathrm{F}}}{T_1}(s^2 - 1)\right) \underset{T_1\to 0}{\longrightarrow} \theta\left(1 - s^2\right) = 0,\ s \gtrsim 1,$$

$$f_0\left[\frac{\varepsilon_{2\mathrm{F}}}{T_2}(\eta s - s_0)^2 - 1\right] \approx f_0\left(2\frac{\varepsilon_{2\mathrm{F}}}{T_2}\eta\delta s\right) \approx f_0(0) = \frac{1}{2},$$

$s_0 \lesssim \eta + 1$.

Now let us consider a case where the impacting Fermi-liquid drop has zero temperature, T_2=0, and the resting drop has temperature T_1. We assume that the following conditions (compare with (49)) are valid:

$$\mathcal{F} \ll 1, \quad s \gtrsim 1, \quad s_0 \lesssim \eta + 1, \quad (\varepsilon_{1\mathrm{F}}/T_1) \gg 1,$$

$$\delta s\,(\varepsilon_{1\mathrm{F}}/T_1) \ll 1, \quad \delta s = s - 1 \ll 1. \tag{56}$$

In this case, by considering the asymptotic estimate (51), the dispersion equation in the main approximation in δs has the form

$$\tilde{\varepsilon}_1(s) \approx 1 + \alpha_1\mathcal{F}(1 - I_0) + \frac{\alpha_2\mathcal{F}}{\eta} -$$
$$-\frac{\alpha_1\mathcal{F}}{2}\ln\frac{4\varepsilon_{1\mathrm{F}}}{T_1} + \frac{\alpha_2\mathcal{F}}{2\eta}\ln\frac{\eta\delta s}{2} = 0. \tag{57}$$

A solution of this equation is given by

$$\delta s = \frac{2}{\eta}\left(4\frac{\varepsilon_{1\mathrm{F}}}{T_1}\right)^{\frac{\eta\alpha_1}{\alpha_2}} \times \exp\left\{-\frac{2\eta Q_2}{\alpha_2\mathcal{F}}\right\}, \tag{58}$$

$$Q_2 \equiv 1 + \alpha_1\mathcal{F}(1 - I_0) + \frac{\alpha_2\mathcal{F}}{\eta}.$$

Here, we should remember that, in accordance with (56), the relation $\delta s\,(\varepsilon_{1\mathrm{F}}/T_1) \ll 1$ must be valid, and a "temperature" condition of the existence of such a solution is determined by the inequality (compare with (53))

$$1 \ll \frac{\varepsilon_{1\mathrm{F}}}{T_1} \ll \exp\left\{\frac{2\eta Q_2}{(\eta\alpha_1 + \alpha_2)\mathcal{F}}\right\}. \tag{59}$$

WE note that, for $s \gtrsim 1$, $s_0 \lesssim \eta + 1$ and $T_2 = 0$, the following formulas are valid:

$$f_0\left(\frac{\varepsilon_{1\mathrm{F}}}{T_1}(s^2 - 1)\right) \approx f_0\left(2\frac{\varepsilon_{1\mathrm{F}}}{T_1}\delta s\right) \approx f_0(0) = \frac{1}{2},$$

$$f_0\left(\frac{\varepsilon_{2\mathrm{F}}}{T_2}\left[(\eta s - s_0)^2 - 1\right]\right) \underset{T_2\to 0}{\longrightarrow} \theta\left(1 - (\eta s - s_0)^2\right),$$

$$f_0\left(\frac{\varepsilon_{2\mathrm{F}}}{T_2}\left[(\eta s - s_0)^2 - 1\right]\right) \underset{T_2\to 0}{\longrightarrow} 1.$$

Then the expression for $\tilde{\varepsilon}_2$, determining the damping or increase of zero-sound oscillations in accordance with (31), can be put in the form

$$\tilde{\varepsilon}_2 = \frac{\pi}{4\eta}\mathcal{F}\left\{\eta\alpha_1 - 2\alpha_2\right\}. \tag{60}$$

We note that the quantity $\{\partial\tilde{\varepsilon}_1/\partial\omega\}_{\omega=\omega_0}$ in compliance with (57) is given by the formula

$$\left\{\frac{\partial\tilde{\varepsilon}_1}{\partial\omega}\right\}_{\omega=\omega_0} \approx \frac{1}{\omega_0}\frac{\alpha_2\mathcal{F}}{2\eta}\frac{1}{\delta s}, \tag{61}$$

where δs is determined by (58). Substituting (60) and (61) into (31), we have

$$\frac{\gamma_k}{\omega_0} \approx \pi\frac{\eta\alpha_1 - 2\alpha_2}{\eta\alpha_2}\left(4\frac{\varepsilon_{1\mathrm{F}}}{T_1}\right)^{\frac{\eta\alpha_1}{\alpha_2}} \times$$
$$\times \exp\left\{-\frac{2}{\alpha_2\mathcal{F}}\left[1 + \alpha_1\mathcal{F}(1 - I_0) + \frac{\alpha_2\mathcal{F}}{\eta}\right]\right\}, \tag{62}$$

$\omega_0 = kv_{1\mathrm{F}}$.

It is easy to see that formula (62) determines, in contrast to (55) which is correct when $T_1 = 0, T_2 \neq 0$, the damping coefficient (when $\eta\alpha_1 > 2\alpha_2$) or a growth (when $\eta\alpha_1 < 2\alpha_2$) of zero-sound oscillations. Such a non-invariance of (55) and (62) relative to the interchange of the impacting and resting drops with the corresponding change of their characteristics, $\varepsilon_{1\mathrm{F}} \leftrightarrow \varepsilon_{2\mathrm{F}}$, $T_1 \leftrightarrow T_2$, (in other words, the absence of the Galilean invariance in the system of collided drops) is not paradoxical in any sense. This is true because, within the developed Fermi-liquid model, the interaction of quasiparticles is not invariant with respect to the Galilean transformations (see (2) and (3)).

Therefore, the relation $\eta\alpha_1 < 2\alpha_2$, which can be written as

$$\alpha_1\sqrt{\varepsilon_{1\mathrm{F}}} < 2\alpha_2\sqrt{\varepsilon_{2\mathrm{F}}}, \tag{63}$$

along with (56) and (59) determines the conditions of the existence of weakly increasing oscillations with increment γ_k in the system:

$$\frac{\gamma_k}{\omega_0} \approx \pi\frac{\eta\alpha_1 - 2\alpha_2}{\eta\alpha_2}\left(4\frac{\varepsilon_{1\mathrm{F}}}{T_1}\right)^{\dfrac{\eta\alpha_1}{\alpha_2}} \times$$
$$\times \exp\left\{-\frac{2\eta}{\alpha_2\mathcal{F}}\left[1 + \alpha_1\mathcal{F}(1 - I_0) + \frac{\alpha_2\mathcal{F}}{\eta}\right]\right\}. \tag{64}$$

In order to finish the study of the temperature impact on the instability development in the system of two collided Fermi-liquid drops, we want to obtain an expression for the increment, when both the resting and impacting drops have comparable ($T_1 \sim T_2$) nonzero temperatures. Moreover,

$$\mathcal{F} \ll 1,\ s \gtrsim 1,\ s_0 \lesssim \eta + 1,\ (\varepsilon_{1\mathrm{F}}/T_1) \gg 1,\ (\varepsilon_{2\mathrm{F}}/T_2) \gg 1,$$
$$\delta s\,(\varepsilon_{1\mathrm{F}}/T_1) \ll 1, \quad \delta s\,(\varepsilon_{2\mathrm{F}}/T_2) \ll 1, \quad \delta s = s - 1 \ll 1. \tag{65}$$

The dispersion equation (46), in the main approximation with regard for conditions (65), has the form

$$\tilde{\varepsilon}_1 = B\delta s^2 - A\delta s + 1 - A = 0, \tag{66}$$

where

$$A \equiv \frac{\mathcal{F}}{2}\left\{\alpha_1 \ln\frac{\varepsilon_{1\mathrm{F}}}{T_1} + \alpha_2 \ln\frac{\varepsilon_{2\mathrm{F}}}{T_2}\right\},$$

$$B \equiv 4\mathcal{F}I_1\left\{\alpha_1\left(\frac{\varepsilon_{1\mathrm{F}}}{T_1}\right)^2 + \eta\alpha_2\left(\frac{\varepsilon_{2\mathrm{F}}}{T_2}\right)^2\right\} \tag{67}$$

and the coefficient I_1 is determined by (51).

Because $(A/B) \ll 1$, the solution of Eq. (66) will be unambiguous and positive when $A \geq 1$:

$$\delta s = \frac{A}{B} + \sqrt{\frac{A^2}{4B^2} + \frac{A-1}{B}}, \quad A \geq 1. \tag{68}$$

Focusing on the explicit form of the coefficients A and B, it is easy to prove that the conditions

$$\delta s\,(\varepsilon_{1\mathrm{F}}/T_1) \ll 1, \quad \delta s\,(\varepsilon_{2\mathrm{F}}/T_2) \ll 1$$

are automatically fulfilled.

In view of (31), (33), we note that, in accordance with (47), (65), and (66), the following formulas hold:

$$\tilde{\varepsilon}_2 \approx \frac{\pi}{4\eta}\mathcal{F}\{\eta\alpha_1 - \alpha_2\},$$

$$\left\{\frac{\partial\tilde{\varepsilon}_1}{\partial\omega}\right\}_{\omega=\omega_0} \approx \frac{1}{\omega_0}\sqrt{A^2 + 4B(A-1)}.$$

Therefore, we get the following expression for the coefficient γ_k:

$$\frac{\gamma_k}{\omega_0} \approx \frac{\pi}{4\eta}\mathcal{F}\frac{\eta\alpha_1 - \alpha_2}{\sqrt{A^2 + 4B(A-1)}} \ll 1, \quad \omega_0 = kv_{1\mathrm{F}}. \tag{69}$$

It is clear that, when $\eta\alpha_1 < \alpha_2$ or

$$\alpha_1\sqrt{\varepsilon_{1\mathrm{F}}} < \alpha_2\sqrt{\varepsilon_{2\mathrm{F}}}, \tag{70}$$

the quantity γ_k represents the instability increment of zero-sound oscillations.

Expression (69) takes a simple form

$$\frac{\gamma_k}{\omega_0} \approx \frac{\pi}{4\eta}\mathcal{F}\frac{\eta\alpha_1 - \alpha_2}{\alpha_1 \ln(\varepsilon_{1\mathrm{F}}/T_1) + \alpha_2 \ln(\varepsilon_{2\mathrm{F}}/T_2)} \tag{71}$$

or

$$\frac{\gamma_k}{\omega_0} \approx \frac{\pi}{4\eta}\mathcal{F}(\eta\alpha_1 - \alpha_2), \quad \omega_0 = kv_{1\mathrm{F}},$$

when $A \approx 1$. That is (see (67)),

$$\alpha_1 \ln(\varepsilon_{1\mathrm{F}}/T_1) + \alpha_2 \ln(\varepsilon_{2\mathrm{F}}/T_2) \approx \frac{2}{\mathcal{F}}. \tag{72}$$

In a more general case, i.e. where $A \gg 1$ and $A \ll B$ (see (67)), formula (69) becomes simpler, but it is still more complicated as compared with (71):

$$\frac{\gamma_k}{\omega_0} \approx \frac{\pi(\eta\alpha_1 - \alpha_2)}{32\eta I_1}\times$$

$$\times\left\{\frac{\mathcal{F}}{2}\left[\alpha_1 \ln(\varepsilon_{1\mathrm{F}}/T_1) + \alpha_2 \ln(\varepsilon_{2\mathrm{F}}/T_2)\right] - 1\right\}^{-1/2}\times$$

$$\times\left\{\alpha_1\left(\frac{\varepsilon_{1\mathrm{F}}}{T_1}\right)^2 + \eta\alpha_2\left(\frac{\varepsilon_{2\mathrm{F}}}{T_2}\right)^2\right\}^{-1}. \tag{73}$$

Expressions for increments (71) and (73) look even simpler, if we consider the collision of identical nuclei having the same temperatures (see formulas (75) and (76) below). The comparison of expressions (38) and (42) with (55), (62), (71), and (73), by taking relations (63) and (70 into account, allows us to draw conclusion that, for the colliding Fermi-liquid drops with nonzero temperatures, the instability increments are higher as compared to those for the colliding drops with zero temperature. In particular, when the colliding Fermi-liquid drops have comparable temperatures $T_1 \sim T_2$ ($(T_1/\varepsilon_{1\mathrm{F}}) \ll 1$, $(T_2/\varepsilon_{2\mathrm{F}}) \ll 1$) and condition (72) is fulfilled, the increase of zero-sound oscillations has a power-like dependence on the Landau amplitude $\mathcal{F} \ll 1$. Whereas, for the other cases considered here (except the one leading to (73)), the degree of the increment smallness is given by exponential multipliers of the type $\exp(-\lambda/\mathcal{F})$, $\mathcal{F} \ll 1$, where λ is a certain constant. We recall that, in our consideration, the collision of two fast (but not relativistic) nuclei is simulated by the collision of two Fermi-liquid drops. Then the given comment means that the instabilities related to the increase of zero-sound oscillations in the system of two colliding excited nuclei (i.e., nuclei with nonzero temperatures) must develop much more intensively as compared with the instabilities in the system of colliding unexcited nuclei.

6. Discussion of Results

Up to now, we have studied the case of small Landau amplitudes $\mathcal{F} \ll 1$. As was proved, the dispersion equations of zero-sound oscillations in this case allow an analytic solution in perturbation theory in small $\mathcal{F}$. The analytic solutions of dispersion equations can be obtained also in the case of large Landau amplitudes, $\mathcal{F} \gg 1$. However, if we consider these solutions from the viewpoint of the problem of the instability development in the system of collided non-relativistic nuclei, it would be easy to prove that a solution of the dispersion equations in perturbation theory in a small parameter $1/\mathcal{F}$ will lead us beyond the non-relativistic approximation which was one of the main assumptions in the present work.

Naturally, it is easy to prove that the solutions of the dispersion equation (30) are within the range of large $s, s \gg 1$ (note at once that, when $s \gg 1$, it is possible to neglect the temperature-dependent components in Eq. (46)). But, in accordance with (31), (33), and (47), the instabilities appear if the condition $s_0 > \eta s$ or

$$\sqrt{\varepsilon/\varepsilon_\mathrm{F}}\cos\alpha > s \gg 1 \tag{74}$$

is satisfied, where ε is the kinetic energy per one nucleon in the impacting drop. In inequality (74), we took into account that, for heavy nuclei, $\varepsilon_{1\mathrm{F}} \approx \varepsilon_{2\mathrm{F}} = \varepsilon_\mathrm{F}$, i.e., $\eta \approx 1$. Remembering that, for the heavy nuclei, $\varepsilon_\mathrm{F} \approx 36$ MeV and the non-relativistic approximation requires the nucleon kinetic energy to be small compared to the rest energy $mc^2 \sim 1$ GeV, ($\varepsilon/mc^2 \ll 1$), it is easy to see that relation (74) cannot be satisfied under the non-relativistic approximation. Therefore, we do not consider the case of large Landau amplitudes in this work.

Here, we have restricted ourselves by consideration of positive Landau amplitudes. When the Landau amplitudes are negative, the dispersion equation of zero-sound oscillations can be solved only by numerical methods. In this case, the real and imaginary parts of the frequency are comparable with each other, and the analytic methods based on perturbation theory become inapplicable. The numerical solution of the dispersion equations for the negative Landau amplitudes is not included into the present work which is aimed at the demonstration of a basic possibility for the instability development in the system of colliding heavy nuclei, which is induced by the propagation and increase of zero-sound oscillations.

Formulas (71), (73) can be used as a proof of the possible development of just such instabilities in the system of heavy colliding nuclei. When two identical nuclei collide ($\varepsilon_{1\mathrm{F}} = \varepsilon_{2\mathrm{F}} = \varepsilon_{\mathrm{F}}$, $\eta = 1$, $T_1 \approx T_2 = T$), formula (71) takes the form

$$\frac{\gamma_k}{\omega_0} \approx \frac{\pi}{2}\frac{\alpha_1 - \alpha_2}{\alpha_1 + \alpha_2}\frac{1}{\ln(\varepsilon_{\mathrm{F}}/T)}, \quad \omega_0 = kv_{\mathrm{F}}. \tag{75}$$

In this case, expression (73) is also significantly simplified:

$$\frac{\gamma_k}{\omega_0} \approx \frac{\pi}{32 I_1}\left(\frac{\alpha_1 - \alpha_2}{\alpha_1 + \alpha_2}\right)(\varepsilon_{\mathrm{F}}/T)^{-2} \times$$

$$\times \left\{(1/2)\,\mathcal{F}(\alpha_1 + \alpha_2)\ln(\varepsilon_{\mathrm{F}}/T) - 1\right\}^{-1/2}, \tag{76}$$

As is known, the temperature T can be related to the excitation energy of a nucleus U by the approximate formula (in this connection, see, for example, [19–22])

$$T \approx \sqrt{aU}, \tag{77}$$

where a is a certain constant which should be determined from experimental data (see, in particular, [19–21]). For example, in accordance with [19] for a nucleus with the mass number $A = 115, a \approx 1/8\,\mathrm{MeV}$; $A = 181, a \approx 1/10\,\mathrm{MeV}$. Note that, within the framework of the Fermi-gas model of nuclei, the quantity a can be found from formula [22]

$$a = \frac{4\varepsilon_{\mathrm{F}}}{A\pi^2}.$$

Following this formula for $\varepsilon_{\mathrm{F}} \approx 36$ MeV and $A = 115, A = 181$, we have $a \approx 0.124\,\mathrm{MeV}$ and $a \approx 0.0795$ MeV, respectively.

If we consider that the instabilities, which are associated with an increase of the amplitudes of zero-sound oscillations, are responsible for the fragmentation, then the specific time of fragmentation τ_f must be of the order of the value reciprocal to the increment γ_k,

$$\tau_f \sim 1/\gamma_k. \tag{78}$$

Thus, if we could experimentally measure the fragmentation time of nuclei τ_f, then the specific dependence of the time on the excitation energy U would be governed by the formula

$$\tau_f \sim C \ln\left(\frac{\varepsilon_{\mathrm{F}}}{\sqrt{aU}}\right), \quad C = \omega_0 \frac{2}{\pi}\frac{(\alpha_1 + \alpha_2)}{(\alpha_1 - \alpha_2)}, \tag{79}$$

when relation (72) holds, and by the formula

$$\tau_f \sim C\left(\frac{\varepsilon_{\mathrm{F}}}{\sqrt{aU}}\right)^2 \left\{\mathcal{F}(\alpha_1 + \alpha_2)\ln\left(\frac{\varepsilon_{\mathrm{F}}}{\sqrt{aU}}\right) - 2\right\}, \tag{80}$$

when condition $A \gg 1$ leading to the validity of (73). According to (75)—(78), this would show that the proposed mechanism of the instability development in the system of colliding heavy nuclei works. In addition, if such instability development mechanism is realized, it can result in the nucleus fragmentation. Consequently, according to (72), (79), and (80), we could estimate both the frequencies of oscillations of the developed instabilities and the Landau amplitudes for the nuclear matter formed as a result of the collision. It is should be noted that the Landau amplitudes are parameters of the Fermi-liquid approach of the present work and, generally speaking, they don't necessarily coincide with the Landau amplitudes for certain heavy nuclei, whose values are given in some published papers (see, for example, [13, 14]).

However, if we make an assumption that values of the Landau amplitudes in the nuclear matter, formed as a result of the collision of two fast heavy nuclei, are close to the corresponding values of the amplitudes of certain nuclei, then we can estimate the increments for specific cases and provide suggestions about the range of using the results of the present work. Work [14] gives grounds to make a conclusion that the value of the Landau amplitude $F_0^{(0)}$ determining oscillations of the nucleus density (see (6)) is within the range of –0.25 and 0.5. These values of the Landau amplitudes from the positive part of this range, when $\varepsilon_{1\mathrm{F}} \approx \varepsilon_{2\mathrm{F}} = 36$ MeV, are fully within the frame of the assumptions made in present work.

The values of the amplitudes $F_0^{(s)}$ and $F_0^{(i)}$ (see (6)) determining oscillations of the spin and isospin densities in the nucleus (see [13, 14]) are also in compliance with the assumptions of our work:

$$F_0^{(s)} \approx 0.27 \div 0.35, F_0^{(i)} \approx 0.59 \div 0.72.$$

As for the negative values of the amplitude F_0 and values of the amplitude $F_0^{(si)}$ (see [13, 14]),

$$F_0^{(si)} \approx 1.26 \div 1.36,$$

determining zero-sound oscillations of the spin-isospin density, these cases require, as we have already mentioned, numerical calculations. This is the reason why they are not included into this work.

Finally, let us study the question about peculiarities in the behavior of a non-equilibrium distribution function of two colliding nucleon drops in the momentum space. This will allow us to determine a direction of nucleus fragment release as a result of the instability development in such a system with respect to the velocity of the moving drop.

It follows from (17) that a deviation of the distribution function $\tilde{\delta f}(\mathbf{p},\mathbf{k},\omega)$ from the equilibrium distribution function is determined by the equation

$$\tilde{\delta f}(\mathbf{p},\mathbf{k},\omega) = -F\left\{\frac{\alpha_1}{\omega-\mathbf{kv}+io}\mathbf{k}\frac{\partial f_0^{(1)}(\mathbf{p})}{\partial \mathbf{p}}+\right.$$

$$\left.+\frac{\alpha_2}{\omega-\mathbf{kv}+io}\mathbf{k}\frac{\partial f_0^{(2)}(\mathbf{p}-m\mathbf{u})}{\partial \mathbf{p}}\right\}\times$$

$$\times\int d\tau'\tilde{\delta f}\left(\mathbf{p}',\mathbf{k},\omega\right)+\delta A(\mathbf{p},\mathbf{k})\,\delta(\omega-\mathbf{vk}), \tag{81}$$

where the derivatives $\dfrac{\partial f_0^{(1)}(\mathbf{p})}{\partial \mathbf{p}}$, $\dfrac{\partial f_0^{(2)}(\mathbf{p}-m\mathbf{u})}{\partial \mathbf{p}}$ are given by (16). According to (4) and (81), the non-equilibrium distribution function can have maximum, if the following conditions are met:

$$(\omega-\mathbf{kv})|_{\mathbf{v}=\mathbf{v}_{1\mathrm{F}}}=0,$$
$$(\omega-\mathbf{ku}-\mathbf{k}(\mathbf{v}-\mathbf{u}))|_{|\mathbf{v}-\mathbf{u}|=v_{2\mathrm{F}}}=0, \quad \mathbf{v}=\mathbf{p}/m.$$

Since $\omega\approx\omega_0=skv_{1\mathrm{F}}$, these conditions can be written as

$$s-\cos\theta=0, \quad \eta s-s_0-\cos\beta=0, \tag{82}$$

where θ is an angle between directions of the vectors $\mathbf{k}$ and $\mathbf{v}$, and β is an angle between directions of the vector $\mathbf{k}$ and the vector $\mathbf{v}-\mathbf{u}$. The first of these conditions is not fulfilled because $s>1$. As we have previously mentioned, a damping or increase of the zero-sound oscillations can take place only when $|\eta s-s_0|<1$. This means that the second condition in (82) can be fulfilled. The maximum value of the increment γ_k is reached when $\eta s-s_0\sim-1$ (see (40)). Under the condition $\varepsilon_{1\mathrm{F}}\approx\varepsilon_{2\mathrm{F}}=\varepsilon_\mathrm{F}(\eta=1)$ and $s\gtrsim 1$, it corresponds to the condition

$$s_0=\frac{u}{v_\mathrm{F}}\cos\alpha\approx 2 \tag{83}$$

(α is, as before, an angle between directions of the vectors $\mathbf{u}$ and $\mathbf{k}$) and, hence,

$$u\geq 2v_\mathrm{F}. \tag{84}$$

The relation $\eta s-s_0\sim-1$ in accordance with (82) means that the vectors $\mathbf{k}$ and $\mathbf{v}-\mathbf{u}$ are antiparallel,

$$\mathbf{v}-\mathbf{u}=-|\mathbf{v}-\mathbf{u}|\frac{\mathbf{k}}{k},$$

or, bearing in mind that $|\mathbf{v}-\mathbf{u}|=v_\mathrm{F}$,

$$\mathbf{v}-\mathbf{u}=-v_\mathrm{F}\frac{\mathbf{k}}{k}.$$

Multiplying scalarwise this equation by vector $\mathbf{v}-\mathbf{u}$ and taking into account (83), we obtain the following relation for the cosine of the angle θ_0 between directions of the vectors $\mathbf{v}$ and $\mathbf{u}$

$$\cos\theta_0=\frac{u^2-2v_\mathrm{F}^2}{vu},$$

Hence, remembering that the modulus of the vector v is within the region of values of $u-v_\mathrm{F}$ and $u+v_\mathrm{F}$ and that, by (84),

$$\frac{u^2-2v_\mathrm{F}^2}{u(u+v_\mathrm{F})}<1, \quad \frac{u^2-2v_\mathrm{F}^2}{u(u-v_\mathrm{F})}\geq 1,$$

we have

$$\frac{u^2-2v_\mathrm{F}^2}{u(u+v_\mathrm{F})}\leq\cos\theta_0\leq 1.$$

In view of inequality (84), the last relation can be written as

$$\frac{1}{3}\leq\cos\theta_0\leq 1. \tag{85}$$

Thus, according to the results of this paper in the case where inequality (84) is true, the jets of the nuclear matter at the collision of heavy nuclei should be expected along the directions defined by condition (85). In the case where inequality (84) is not satisfied, the angular distribution of the outgoing matter should be close to an isotropic one.

Let us come back to the problem of the collisionless approximation for the description of the dynamics of the nuclear matter formed by the collision of heavy nuclei. The main condition of its applicability (1) from the point of view of the results of this paper is the inequality $\omega_0\tau_r\gg 1$, where $\omega_0=skv_\mathrm{F}$ and τ_r is the relaxation time. It is necessary also that the characteristic times of the instabilities' development $\tau_f \sim 1/\gamma$ (where γ are the increments found in this paper) should be smaller or have the same magnitude as the relaxation time $\frac{1}{\gamma} \lesssim \tau_r$. The last inequality can be written as $\frac{\gamma}{\omega_0}\omega_0\tau_r \gtrsim 1$. Thus, the inequalities

$$\omega_0\tau_r \gg 1, \quad \frac{\gamma}{\omega_0}\omega_0\tau_r \gtrsim 1.$$

$$\omega_0 \gg \gamma \gtrsim 1/\tau_r, \quad t \ll 1/\gamma \lesssim \tau_r.$$

are the applicability conditions of the collisionless approximation for the description of the initial stage of development of the instabilities in the nuclear matter formed by the collision of heavy nuclei. Let us remark that the kinetic equation without taking into account collisions between particles had been used by some authors for the description of the dynamics of the matter formed by the collision of heavy nuclei (see, in this case, [1, 2, 23], and references therein). As a rule, this equation was written in the mean field approximation. For example, the spinodal fission in the expanding nucleon Fermi-liquid was investigated in [23] on the basis of such kinetic equation. However, as is known, the average field approach needs considerable numerical calculations. The usage of the Fermi-liquid approach with the phenomenological parameters of interaction (Landau amplitudes) for describing the dynamics of the nuclear matter allows one, in many cases, to use the analytic treatment without numerical methods.

As was repeatedly noted in this paper, the application of the Fermi-liquid approach is justified while describing the properties of heavy nuclei. For this reason, the confirmations of the mechanism of the origin of instabilities in the nuclear matter suggested in this paper can be expected in experiments at the collisions of heavy fast (but non-relativistic) nuclei. These might be the reactions Gd+U or Xe+Sn (INDRA, [2, 12]) at the energies of the incoming nucleus of more than 145 MeV per nucleon.

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ФЕРМІ-РІДИННИЙ ПІДХІД ДЛЯ ОПИСУ ПОЧАТКОВОЇ СТАДІЇ ФРАГМЕНТАЦІЇ ПРИ ЗІТКНЕННЯХ ВАЖКИХ ЯДЕР

А.П. Івашин, С.В. Пелетминський, Ю.В. Слюсаренко

Р е з ю м е

Запропоновано механізм розвитку початкової стадії нестійкості, що може приводити до фрагментації ядерної матерії, яка виникає внаслідок зіткнень нерелятивістських важких ядер. Зіткнення важких ядер моделюється як зіткнення двох необмежених фермі-рідинних "крапель". Походження нестійкості в такій системі пов'язується із поширенням наростаючих коливань в ядерній матерії. Ці коливання можуть існувати у фермі-рідині, що перебуває у стані спокою: модифікований нульовий звук Ландау, модифіковані спінові та ізоспінові хвилі, комбінація цих хвиль. Указані нестійкості аналогічні пучковій нестійкості у звичайній електронній плазмі. Наведено аналіз отриманих інкрементів коливань, який можна використати для експериментального підтвердження запропонованого механізму фрагментації при ядерних зіткненнях. Конкретизуються напрямки, уздовж яких очікуються "струмені" ядерної матерії.