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Control of Strongly Nonequilibrium Coherently Correlated States and Superconducting Transition Temperature

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Abstract: Our paper considers the possibility of the emergence and control of non-equilibrium states of a quasi-homogenous condensed medium with energy and particle flows in the phase space, which, first of all, manifest themselves in the explosive development of the asymmetry in the initially symmetric equilibrium system. This symmetry breaking and the appearance of non-equilibrium in the system are controlled by the coherent acceleration of the system. Dependencies of thermodynamic parameters of a strong nonequilibrium system on the indices of disequilibrium in coherently correlated states are given, and the estimates of the dielectric permittivity in a non-equilibrium system and modes of plasma acoustic oscillations are made. An estimate of the superconducting transition temperature under nonequilibrium conditions has been made. It is demonstrated that the superconducting transition temperature can approach the limiting value, corresponding to a quantum with its plasma frequency of the medium.

Keywords: strongly nonequilibrium distributions; coherent acceleration; coherently correlated states; energy flow in the phase space; phase transition temperature

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1. Introduction

Advances towards creating methods for controlling the direction of evolutionary processes are largely based on the development of nonequilibrium thermodynamics, and interest in these issues continues unabated [1–5]. The emerging understanding that evolution under certain conditions can be influenced by very insignificant impacts [6–8] makes it clear that there is growing concern about the possible existence of unconscious, unpredictable controlling influences of random electromagnetic pollutants on the ecology of the external environment and on biological systems. There is a pressing need to solve problems of managing self-organization in various complex systems.

It should be noted that despite huge achievements, Prigogine's nonequilibrium thermodynamics [1] is limited by an approximation of local equilibrium states, the strong non-equilibrium of which lies only in the strong spatial dependence of the parameters of equilibrium distributions (density, temperature, etc.). Naturally, the management of system evolution and phase transitions, in this case, is carried out through variations of these thermodynamic parameters. The peculiarities of strongly nonequilibrium states in quasi-homogeneous systems and the peculiarities of non-locality influence in coherently correlated states of complex systems on the sequence of phase transitions and, therefore, on the evolution of their internal structure are insufficiently studied. The key role of correlations in many-particle systems was first clearly demonstrated in the fundamental works of N. Bogoliubov on kinetic equations and hydrodynamics [9] (the results of which are detailed, for example, in Refs. [9,10]).

A peculiar theory of kinetic equations considering the growth and evolution of structures in a particle system and non-locality of system states was obtained and analyzed by Vlasov et al. [11–13] using Cartan geometry (the physical meaning of Cartan geometry is thoroughly demonstrated in Penrose's work [14]). Vlasov's nonlocal statistical mechanics of a many-particle system and the theory of structure growth are set out in [15–17], where the non-locality of particle states is the main axiom of his evolutionary theory.

To verify the non-locality of physical phenomena, Bell proposed a theorem (see, for example, [18]) with inequalities, which now bear his name. The fundamental importance of this theorem is emphasized in [19]: "What Bell's theorem, together with the experimental results, proves to be impossible (subject to a few caveats we will attend to) is not determinism or hidden variables or realism but locality, in a clear sense. What Bell proved, and what theoretical physics has not yet properly absorbed, is that the physical world itself is non-local." The experimental proof of violations of Bell's inequalities, and therefore the objective non-locality of particle interaction, was awarded the Nobel Prize in 2022 [20]. The official formulation of the Nobel Committee is more specific and reads—"for experiments with entangled photons, establishing violations in Bell's inequality, and for innovations in quantum informatics".

Non-locality in the dynamics and kinetics of particles leads to the fact that the influence on physical processes is not limited to the infinitesimal vicinity of the considered point but requires accounting for the finiteness of the interaction region and considering all orders of differentials of variables describing the system state. There are two qualitative approaches to describing these effects:

- Operator regularization method: Instead of differential operators for describing the evolution of the system, for example, functional, integrodifferential equations ([21–23]) and flows in the phase space are used (see, for example, in the Refs. [24–33]).
- Regularization method for the evolution of systems: instead of the four-dimensional spacetime, for example, the Cartan space is used (see, for example, the works of Vlasov [15–17]).

In the nonlocal theory of Vlasov, the spacetime in which the statistical theory of dynamic systems should be formulated is not the usual four-dimensional Riemann spacetime but the Cartan space. The Cartan space is a combination of points in the four-dimensional Riemann spacetime and tangent surfaces of different orders at each point.

The main variables in the statistical theory are not only coordinates, time, and velocities (momenta) but also accelerations of all orders. In his covariant kinetic equation in the Cartan space (see Ref. [16]), Vlasov obtained the exact power-law solution for non-inertial reference systems, i.e., systems in strong nonequilibrium with coherent acceleration.

Physically equivalent descriptions of non-inertial systems with coherent acceleration are spatially homogeneous systems of particles with phase space flows [24–28]. Exact power-law solutions for kinetic equations of Boltzmann, Landau, Balescu-Lenard [25–27], and kinetic equations in fractal media (see [24]) have been obtained for such physical models.

In all the mentioned cases, strongly nonequilibrium particle distributions in phase space exhibit power-law behavior, which reflects the common physical nature of evolution control—coherent acceleration and phase space flow (see [28]). It is important to note that coherent acceleration and the corresponding inertial forces are associated with the variation of the system's structure, i.e., its binding energy and inertia [28–30].

Energy flows in the system arise either due to changes in the environment or due to internal changes in the structure of the system itself (see [28–30]). Importantly, environmental changes can influence both through the boundary between the designated system and the environment and uniformly throughout the system if the considered system is fully immersed in the environment. Each element of the system experiences the influence of the environment.

Section 2 of this study presents the Vlasov equations on the first tangent bundles after the corresponding averaging over higher accelerations. In these Vlasov equations, we explicitly account for the mentioned dependencies between variations in the system's structures, coherent acceleration of many-particle systems, and particle and energy fluxes, which affect the distributions of charged particles and the characteristics of phase transitions (see for example, Refs. [31,32]). They can lead to a significant change in the phase transition temperature, the estimation of which is provided at the end of this article.

Together with the averaged kinetic equations, we obtain the dynamic transport equations. The dynamic equations in non-inertial reference systems, when averaging over higher orders of acceleration and considering the evolution of the internal structure of the system, can exhibit positive feedback that is consistent with the results obtained in [33]. In the same section, we present the covariant Vlasov kinetic equation with the evolution of the internal structure considered, along with its exact power-law solution expressed through non-extensive generalizations of Tsallis exponentials [34–37]. The obtained non-equilibrium distributions in systems with coherent acceleration are consistent with the general properties of accelerated reference systems [38] and the geometric properties of systems with constraints [39–42].

The processes of system evolution from an equilibrium system to a non-equilibrium quasi-stationary state with energy flows and/or the number of particles P_S in the phase space of the system (non-equilibrium phase transitions) are accompanied by the emergence and explosive power-law growth of the asymmetry of the system's phase space, the structure of the phase space and the corresponding evolution of the order parameter η.

Section 3 of this study presents the fundamental thermodynamic properties of systems in strongly non-equilibrium states close to spatial homogeneity. The physical situation where each element of the system experiences the same influence from its surrounding environment is crucial for the system's evolution and corresponds to the action of a mass force (and the corresponding flow P_S), leading to coherent acceleration of the entire system [28–30]. This mass force can result from changes in the environment's structure or contributions from changes in the internal structure of the system itself, which occur under the influence of mass forces and coherent acceleration. The properties of nonlinearity, nonlocality, and nonextensiveness [34–37] are closely related to the thermodynamic properties of coherently correlated systems (see [43–46]).

The properties of nonlinearity, nonlocality, and nonextensiveness (coherent parameter q) [34] are closely related to the thermodynamic properties of coherently correlated systems (see [47,48]), anisotropy, and the properties of thin films [49]. In [50], the role of local nonequilibrium of electron states in the processes of increasing the temperature of the superconducting transition is shown. The non-equilibrium phase transition of evolution to a non-equilibrium quasi-stationary state is accompanied by a certain dependence between the coherence parameter q, the order parameter η and flows P_S in its phase space.

In Section 4 of this study, the obtained relationships for non-equilibrium states and particle distributions are utilized to estimate the dielectric permeability of the non-equilibrium medium and the parameters of acoustic modes in these non-equilibrium states. At the end of the section, using the "jelly" model [51] and Ginzburg's relations for the temperature of the superconducting transition [52], temperature estimations for the superconducting transition in these non-equilibrium states are derived.

2. On the Vlasov Kinetic Equations in Systems with Varying Internal Structure

The body forces and coherent accelerations of all orders cause the non-locality and evolution of the system (changes in its internal structure, see, for example, [15,16]). The geometric representation of particle dynamics leads to the general Vlasov kinetic equation for the distribution function $f\left(t, \vec{r}, \vec{u}, \vec{a}, \dot{\vec{a}}, \ddot{\vec{a}}, \ldots..\right)$, which has the form of a continuity equation in Cartan space:

$$\frac{\partial f}{\partial t} + div_{\vec{r}}\left(\vec{u} f\right) + div_{\vec{u}}\left(\vec{a} f\right) + div_{\vec{a}}\left(\dot{\vec{a}} f\right) + \ldots. = 0 \tag{1}$$

The divergent character of the kinetic equation corresponds to the free motion of particles in this complete space. From the kinetic Equation (1) for the distribution function $f\left(t, \vec{r}, \vec{u}, \vec{a}, \dot{\vec{a}}, \ddot{\vec{a}}, \ldots..\right)$ in the full space, by successive averaging over higher-order accelerations, one can obtain a set of averaged physical quantities:

$$\rho\left(t,\vec{r}\right) = \int d\vec{u} f\left(t,\vec{r},\vec{u}\right),\ f\left(t,\vec{r},\vec{u}\right) = \int d\dot{\vec{u}} f\left(t,\vec{r},\vec{u},\dot{\vec{u}}\right),\ f\left(t,\vec{r},\vec{u},\dot{\vec{u}}\right) = \int d\ddot{\vec{u}} f\left(t,\vec{r},\vec{u},\dot{\vec{u}},\ddot{\vec{u}}\right), \quad (2)$$

describing a reduced description, a chain of linking equations containing several divergent terms that correspond to the number of tangent bundles over which no averaging has been performed. Together with the sequence of kinetic equations that correspond to the hierarchy of averaging, the dynamic equations for the main average kinematic quantities (the corresponding equations of motion) also follow [10,16].

Below, we present the Vlasov kinetic equations [16,17] considering the variation of the internal structure during evolution [28] at the first three main levels of the averaging hierarchy.

2.1. Averaging over all Tangent Spaces, Averaging over Speeds and accelerations of All Orders

With this averaging, we obtain the continuity equation for the density and the equation of motion in a dissipative medium:

$$\frac{\partial\rho}{\partial t} + div_{\vec{r}}\left(\left\langle\vec{u}\right\rangle f\right) = 0,\ \left\langle\nu_{dis}\vec{u}\right\rangle = \frac{1}{m}\vec{F}\left(t,\vec{r},\rho\right). \quad (3)$$

Vlasov called the resulting particle dynamics equation the Aristotelian equation of motion—the average speed when moving in a dissipative medium is proportional to the force.

2.2. Averaging over Tangent Spaces of Accelerations of all Orders

In this case, we obtain the Vlasov kinetic equation for the distribution function in the phase space (in terms of coordinates and impulses or velocities) and the averaged equations of dynamics, which coincide in inertial reference frames with Newton's dynamic equations in the first tangent bundle:

$$\frac{\partial f_r}{\partial t} + div_{\vec{r}}\left(\vec{u} f_r\right) + div_{\vec{u}}\left(\left\langle\vec{a}\right\rangle f_r\right) = 0;\ \left\langle\frac{d}{dt}\left(m\vec{u}\right)\right\rangle = \vec{F}(t,\vec{r},\vec{u},f_r). \quad (4)$$

This approximation is sufficient in inertial frames of reference, but in non-inertial frames of reference, Equation (4) must include evolutionary processes with a change in the internal structure of the system (entropy). In non-inertial systems, the left-hand side should contain the change in momentum, taking into account viscous dissipative losses $\nu_{dis}\left\langle\vec{u}\right\rangle$ when interacting with the medium and variations in bonds and internal structure (and, therefore, entropy) in the base space and the first tangent bundle takes the form $\left(\frac{1}{m}\frac{d}{dt}m(S)\right)\left\langle\vec{u}\right\rangle$. And the dynamics equation from (4) takes the form:

$$\frac{d}{dt}\langle\vec{u}\rangle + \nu_{eff}\langle\vec{u}\rangle = \frac{1}{m}\vec{F}_m(t,r,\vec{u},\{f_r\})$$

where $\vec{F}_m\left(t,r,\vec{u},\{f_r\}\right)$ is external mass force acting on the system, $\nu_{eff} = \left(\nu_{dis} + \frac{1}{m}\frac{d\,m(S)}{dt}\right)$—effective self-consistent viscosity, and $\nu_{eff}\left\langle\vec{u}\right\rangle$ is the averaged self-consistent force of viscosity between particles and the medium, taking into account the entropy force arising due to the variation of inertia (see, for example, [18–20]). It can be seen that, under certain conditions, bond variations can make the effective self-consistent viscosity ν_{eff} negative [53,54].

Let us now consider the kinetic equation from (4). The mass force $\vec{F}_m\left(t, r, \vec{u}, \{f_r\}\right)$ initiating the evolution process creates an initial coherent acceleration $\vec{a}_{cog}$ for the subsystem in the coherence region with the scale: $l_{cog} = \frac{c^2}{2\,a_{cog}}$ [39–42]. As a result, the subsystem becomes coherent and homogeneous within the region with the specified scale, and variations in the structure and binding energy of the system occur, contributing to the self-consistent viscous force.

We write the entropy forces that modify the Vlasov equation in the form of a collision integral in the divergent form:

$$\begin{gathered}\frac{\partial f\left(\vec{r},\vec{u},t\right)}{\partial t} + div_{\vec{r}}\left(\vec{u}\, f\left(\vec{r}, \vec{u}, t\right)\right) + div_{\vec{u}}\left(\left\langle \dot{\vec{u}} \right\rangle f\left(\vec{r}, \vec{u}, t\right)\right) = div_{\vec{u}}\left(\vec{j}_S\right); \\ \vec{j}_S = \left(-\frac{1}{m}\nabla_r\left(\omega_{eff} S_q\right)\right) f\left(\vec{r}, \vec{u}, t\right);\ \left\langle \dot{\vec{u}} \right\rangle = \frac{1}{m}\vec{F}_B,\ \omega_{eff} \approx \frac{2\pi}{\tau_{eff}}.\end{gathered} \tag{5}$$

The action of the mass entropy force on a system of particles and the acceleration of the system, which is inextricably linked with it, force the system to rebuild its internal structure and, thus, evolve in the tangent bundle of space-time in accordance with the variational principle of dynamic harmonization.

The kinetic Equation (5) for the particles included in the subsystem in the eight-dimensional space of the supporting elements can be represented in the covariant form:

$$\widehat{D}iv_r\left(\vec{u}\, f\right) + div_u\left(\left\langle \frac{\widehat{D}\vec{u}}{d\tau} \right\rangle f\right) = 0 \tag{6}$$

Given that:

$$\widehat{D}iv_r\left(\vec{u}\, f\right) = u^\alpha \widehat{D}_\alpha f + f \widehat{D}_\alpha u^\alpha,\ div_u\left(\left\langle \frac{\widehat{D}\vec{u}}{d\tau} \right\rangle f\right) = \frac{\partial}{\partial u^\alpha}\left(\left\langle \frac{\widehat{D}\vec{u}}{d\tau} \right\rangle^\alpha f\right),$$

$$\widehat{D}_\alpha f = \frac{\partial f}{\partial x^\alpha} - \Gamma^\sigma_{\alpha\gamma} u^\gamma \frac{\partial f}{\partial u^\sigma},\ \Gamma^\sigma_{\alpha\gamma} = \frac{1}{2} g^{\mu\sigma}\left(\frac{\partial g_{\mu\alpha}}{\partial x^\gamma} + \frac{\partial g_{\mu\gamma}}{\partial x^\alpha} - \frac{\partial g_{\alpha\gamma}}{\partial x^\mu}\right),$$

we write (6) by components:

$$u^\alpha \frac{\partial f}{\partial x^\alpha} - \Gamma^\sigma_{\alpha\gamma} u^\gamma \frac{\partial f}{\partial u^\sigma} + \frac{\partial}{\partial u^\alpha}\left(\left\langle \frac{\widehat{D}\vec{u}}{d\tau} \right\rangle^\alpha f + P_s\right) = 0. \tag{7}$$

Let us analyze the solutions of the kinetic equation in an important particular case, when the explicit contribution of the divergence $div_{\vec{u}}(.)$ into the kinetic equations can be ignored, and the covariant equation of quasi-stationary states takes the form:

$$u^\alpha \frac{\partial f}{\partial x^\alpha} - \Gamma^\sigma_{\alpha\gamma} u^\gamma \frac{\partial f}{\partial u^\sigma} = 0. \tag{8}$$

This approximation is valid in two cases:

- when there are no external forces in the system and no flow in the phase space $P_S = 0$ (this case corresponds to complete equilibrium in the system);
- when external forces do not act directly in the system, but the flow of energy, particles or entropy is constant in the phase space $P_S = const$ (this corresponds to a strong deviation from equilibrium).

Let us consider solutions of the kinetic Equation (7) isotropic in the space of velocities and stationary in laboratory time, i.e., Let's pretend that:

$$\partial f/\partial x^0 = 0;\ \partial g_{\alpha\beta}/\partial x^0 = 0,\ g_{0i} = 0,\ i = 1,2,3.$$

We emphasize that this stationarity does not imply independence of proper time. In this approximation:

$$\Gamma^0_{00} = \Gamma^0_{ik} = 0;\ \Gamma^k_{i0} = 0;\ \Gamma^0_{0i} = \frac{1}{2g_{00}}\frac{\partial g_{00}}{\partial x^i};\ \Gamma^i_{00} = -\frac{1}{2}g_{ik}\frac{\partial g_{00}}{\partial x^k}.$$

Separating the variables, we represent the distribution function in the form:

$$f\left(x^\alpha, u^\alpha, u^0, t\right) = f\left(x^\alpha, u^\alpha, u^0\right) = \rho(x^\alpha)\psi\left(u^2\right)\psi_0\left(u^0\right), \tag{9}$$

$$(1-q)\xi^2 = g_{\alpha\beta}\xi^\alpha\xi^\beta,\ \xi^\alpha = \frac{u^\alpha}{\sqrt{-(1-q)g_{oo}}u^0},\ u^2 = u_\alpha u^\alpha. \tag{10}$$

After these relations, from (8) follows the system of equations, which is exactly solved, and the solution can be expressed in terms of the generalized exponential $exp_q(x)$ [34]:

$$\rho\left(x^i\right) = \rho_0 exp_q\left(-\frac{U(x)}{w(q)}\right),\ \psi_0\left(u^0\right) = \psi_0(0)\left(u_0{}^{\frac{1}{1-q_{cr}}}\right)^{q_{cr}},\ \psi\left(\xi^2\right) = \left(exp_q\left(-\xi^2\right)\right)^q,$$

$$exp_q(x) = \begin{cases} 1+(1-q)x > 0,\ (1+(1-q)x)^{\frac{1}{1-q}} \\ 1+(1-q)x \le 0,\ 0 \end{cases} \tag{11}$$

The dependence of the distribution function $F_q(\varepsilon/T) = \sqrt{\varepsilon/T}exp_q(\varepsilon/T)$ on energy and coherence index q is shown in Figures 1 and 2.

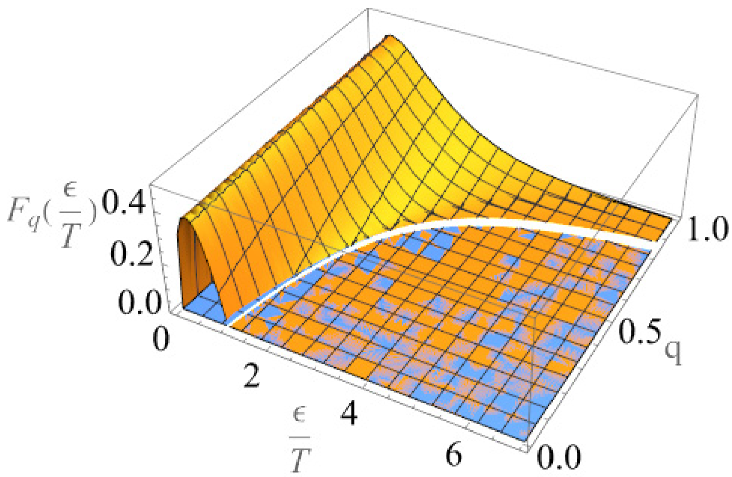

Figure 1. Dependence of the distribution function $F_q(\varepsilon/T)$ on energy ε/T and coherence parameters $q < 1$. With increasing deviation from equilibrium, there is an increase in localization. (The function values are represented in one color, and the zero value is represented in blue. The surface rises above the zero value.)

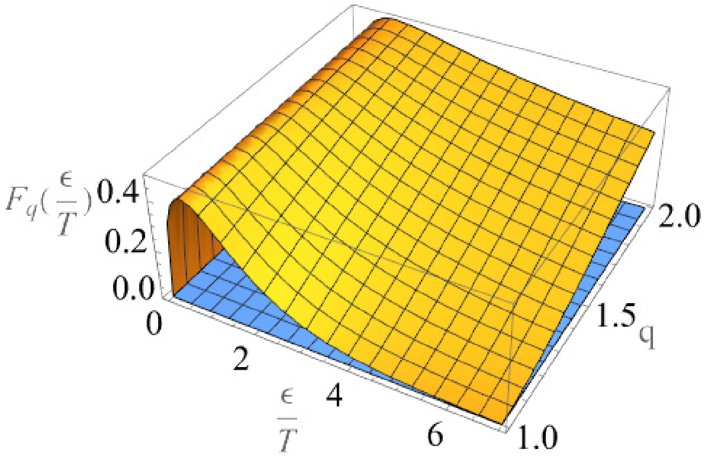

Figure 2. Dependence of the distribution function $F_q(\varepsilon/T)$ on energy ε/T and coherence parameter $q > 1$. With increasing deviation from equilibrium, there is an increase in delocalization. (Non-zero function values are represented in a single color, and a zero value is represented in blue. The surface rises above the zero value.)

These solutions reflect that in the absence of an entropy flow (i.e., at $q = 1$), a homogeneous equilibrium case is realized. The distribution function $\psi(\xi^2)$ for velocities (or energies) passes into the Maxwell distribution function, the distribution function $\rho(x^i)$ into the Boltzmann distribution, and the function $\psi_0(u^0)$ becomes a constant. There is no difference between local and laboratory time. To the extent that entropy flows are present in the system ($q \neq 1$), the distribution functions over energies and coordinates change from an exponential dependence to a quasi-power one.

An analysis of the behavior of the solutions shows that for $q < 1$, the distribution is localized with an increase in the deviation from 1, and for $q > 1$, an increase in the deviation leads to an increase in the delocalization of the distribution function and with it an increase in the characteristic dispersion of physical quantities that depend on momenta and energies.

Exact solutions of the kinetic equations of a function with power asymptotic were studied in [24–27], and the thermodynamic properties of similar functions and states were obtained in many works [5,28–30,34–37]

An analysis of the exact solutions of the kinetic equations and thermodynamic functions in nonequilibrium states leads to the conclusion about the relationship between the coherence parameter q and fluxes P_S in the phase space in accordance with the relation $q \approx \sqrt{1 + P_S}$ [29,30]. In this case, the sign of the flow in the phase space determines the region $q < 1$ and region $q > 1$. The coherence parameter q also can be related to the order parameter $0 \leq \eta \leq 1$ using the relations:

$$q(\eta) = \begin{cases} q_- = 1 - \eta, \ q \leq 1 \\ q_+ = \frac{1}{1-\eta}, \ q > 1 \end{cases} \tag{12}$$

2.3. *Averaging over Tangent Spaces of Higher-Order Accelerations*

With this averaging, we obtain the kinetic equation, taking into account the dynamics of the system in the base space and the first two tangent bundles:

$$\frac{\partial f}{\partial t} + div_{\vec{r}}\left(\vec{u} f\right) + div_{\vec{u}}\left(\vec{a} f\right) + div_{\vec{a}}\left(\left\langle \dot{\vec{a}} \right\rangle f\right) = 0$$

$$\frac{d}{dt}\left\langle \frac{d}{dt}\left(m\vec{u}\right)\right\rangle + \nu_a \left\langle \frac{d}{dt}\left(m\vec{u}\right)\right\rangle = \vec{F}_1\left(t, \vec{r}, \vec{u}, \dot{\vec{u}}, f\right) \tag{13}$$

In a non-inertial frame of reference, the equation of dynamics (13) for the average value of acceleration, considering the variation of inertia, takes the form:

$$\frac{d^2\left\langle \vec{u} \right\rangle}{dt^2} + (2\sigma_S + \nu_a)\frac{d\left\langle \vec{u} \right\rangle}{dt} + \left(\frac{d\sigma_S}{dt} + \sigma_S^2 + \nu_a\sigma_S\right)\left\langle \vec{u} \right\rangle = \dot{\vec{a}}_{cog}\left(t, \vec{r}, \vec{u}, \dot{\vec{u}}, f\right)$$

$$\dot{\vec{a}}_{cog}\left(t, \vec{r}, \vec{u}, \dot{\vec{u}}, f\right) = \frac{1}{m}\vec{F}_1\left(t, \vec{r}, \vec{u}, \dot{\vec{u}}, f\right),\ \sigma_S = \frac{1}{m}\frac{dm}{dt}. \tag{14}$$

Within the coherence scale, the system can be considered homogeneous with high accuracy, and the term $div_{\vec{r}}\left(\vec{u}f\right)$ in the kinetic equation can be neglected. In this case, of course, the main evolution now occurs in tangent bundles—spaces of velocities (energy space) and accelerations, which form a layered energy phase space and satisfy Equations (11) and (12). In Equation (12) $\vec{F}_1\left(t, \vec{r}, \vec{u}, \dot{\vec{u}}, f\right)$ is the external force averaged over the derivatives of accelerations above the first and acting on the system in the energy phase space with coordinates $\left(\vec{u}, \vec{a}\right)$, and $\dot{\vec{a}}_{cog}\left(t, \vec{r}, \vec{u}, \dot{\vec{u}}, f\right)$ the corresponding value of the coherent velocity acceleration (averaged coherent value of the third derivative of the coordinate of the center of gravity of the excited system).

The value $\alpha\left\langle \vec{a}_p \right\rangle$ is the average viscosity force between the particles and the medium in the energy phase space. The viscous force $\nu_a\left\langle \vec{a}_p \right\rangle$ for a system of charged particles is the force of radiative friction [55]:

$$\nu_a\left\langle \vec{a}_p \right\rangle = -\frac{1}{m}\vec{F}_{rad},\ \vec{F}_{rad} = \frac{2}{3}\frac{q_e^2}{c^3}\frac{d^2\vec{u}}{dt^2}.$$

In the case of harmonic excitation with a frequency, the radiative friction force can be represented as a viscous friction force in the velocity space (energy space):

$$\vec{F}_{rad} = \frac{2}{3}\frac{q_e^2}{c^3}\frac{d^3\vec{r}}{dt^3} = D_\tau u,\ D_\tau = \frac{2}{3}\tau_e m_e \omega^2,\ \tau_e = \frac{q_e^2}{m_e c^3} = \frac{r_e}{c}. \tag{15}$$

Here we introduced the coefficient of friction D_τ and the characteristic time τ_e, i.e., the time it takes the light to travel through the effective radius of an electron. In accordance with the equations of motion (13), positive feedback arises, and the coherent acceleration grows exponentially when the threshold is exceeded (determined by the dispersion in the acceleration space and the temperature of the chaotic component of the dynamics of the system of particles and the medium). The equations of motion obtained from the Vlasov kinetic equations agree with the relations for fractal media obtained in [34].

When comparing the equation of dynamics with allowance for radiative friction and the equation for the evolution of a system under the action of a body force, one can see that the forces of radiative friction manifest themselves as entropy forces, which are determined by a change in some structure.

Coherent acceleration of a system of particles leads to the anisotropy of the system—its properties and distribution functions are different in the direction of coherent acceleration and on surfaces orthogonal to acceleration and provide qualitatively different thermodynamic properties.

3. On the Thermodynamics of Coherent-Correlated States of Complex Systems

The degree of order and the probability of the state of the system is characterized by the number of its possible states Ω, the optimization of which implies the form of the nonequilibrium distribution function (see, for example, [24]).

Impulse actions on the surface of condensed media (see [28,29]) can single out spatially thin regions with high correlations, in which the set of particles have the properties of strongly nonequilibrium spatially homogeneous systems with flows in the phase space. The flows in the phase space of the system determine partitioning into the corresponding physically infinitesimal elements of the system and rearrangement of the phase space. The properties of the partitioning of the phase space (and, consequently, the thermodynamics of the system) depend significantly on correlations in the phase space $0 \leq k_{rp} \leq 1$. In this case, to estimate the magnitude of a physically infinitesimal volume, it is necessary to use the generalized instead of the Heisenberg uncertainty relations Schrödinger-Robertson ratio: $\Delta x \Delta p_x \geq \frac{1}{2}\hbar_{eff}$, where $\hbar_{eff} = \hbar\frac{1}{\sqrt{1-k_{rp}^2}}$ (see [43–46]). Moreover, it should be considered that correlations k_{rp} in the phase space (its internal fractal structure and coherent acceleration) are related to the order parameter η in a power-law manner, for example, $k_{rp} \approx \eta^{\alpha_{str}}$ with an index α_{str}, which depends on the structure of the phase space.

When describing such nonequilibrium systems, it becomes necessary to use the concept of a coherently correlated state [27–29], especially for anisotropic systems such as shells and thin layers of condensed matter. We note that size quantization modes and nonequilibrium distribution functions with several components localized in different energy regions can appear in thin layers of condensed matter. Quantization modes and acoustic plasma oscillations [49] can arise in such systems, consistent with coherently correlated states and significantly affect superconducting transitions.

Due to nonequilibrium properties, which are consistent with coherently correlated states, these states have specific thermodynamic properties and acoustic plasma oscillations are initialized. The role of locally nonequilibrium functions in increasing the temperature of the superconducting transition was studied in [50].

Thermodynamic relations in a coherently correlated state can be obtained using Lagrange multipliers in optimizing a non-extensive generalization of entropy for a fixed number of particles and energy of the nonequilibrium state.

Defining the distribution functions directly using the thermodynamic formalism in the corresponding tangent bundles is convenient. The universal distributions obtained above from nonlocal kinetic equations (see also [29,30]) are realized in self-similar systems on different scales where the phase space has a fractal structure. To construct the statistics of such systems, we apply the important approach of K. Tsallis [34]. He proposed, when determining the entropy and distribution of states, to deform the logarithmic function in such a way that $\ln_q(x) = \frac{x^{1-q}-1}{1-q}$, for large values of the energy of the states, the probability of their realization would decrease not exponentially quickly but in a power-law manner (Pareto's law). The function $\ln_q(x)$ is the inverse function to the function $exp_q(x)$ in terms of which distributions are expressed in coherently accelerated systems.

Let the system with the number of particles N have a subsystem with distinguished coherence properties and the number of particles N_{cog}. These numbers of particles naturally determine the order parameters by the relationship: $\eta = \frac{N_{cog}}{N}$, $N = N_{cog} + N_g$.

The most probable distribution functions are determined by the entropy extremum. To determine the necessary extrema of entropy, considering the restrictions, we use the method of Lagrange multipliers. The required entropy optimum for the distribution function $f_i(\varepsilon_i)$, taking into account the restrictions on the number of particles and energy: $\sum_{i=1}^{w} \varepsilon_i f_i = W$, $\sum_{i=1}^{w} f_i = \eta$ is found as the unconditional extremum of the function $\widetilde{S}(f_i, \alpha, \beta)$, which is the entropy averaged with the function f_i: $\sum_i S_i f_i = -\sum_i q\frac{(f_i{}^{1-q}-1)}{1-q} f_i$, supplemented by a linear combination of functions that characterize the constraints:

$$\widetilde{S}(f_i, \alpha, \beta) = -\frac{q}{1-q}\sum_i \left(f_i{}^{2-q} - f_i\right) - \alpha\left(\sum_{i=1}^{w} f_i - \eta\right) - \beta\left(\sum_{i=1}^{w} \varepsilon_i f_i - W\right).$$

The condition for the unconditional extremum of the function $\widetilde{S}(f_i, \alpha, \beta)$ in all variables f_i, α, β gives the system of equations:

$$-\frac{q(2-q)}{1-q} f_i^{1-q} + \frac{q}{1-q} - \alpha - \beta\varepsilon_i = 0 \quad (16)$$

$$\left(\sum_{i=1}^{w} f_i - \eta\right) = 0, \; \left(\sum_{i=1}^{w} \varepsilon_i f_i - W\right) = 0 \quad (17)$$

from which follows the expression for the optimal function:

$$f_i = A_q exp_q\left(-\frac{\varepsilon_i}{T_q}\right) \quad (18)$$

$$T_q(\alpha, \beta) = \frac{(q + (q-1)\alpha)}{\beta}, \; A_q(\alpha, \beta) = \left(\frac{1}{2-q}\left(1 + \frac{(q-1)}{q}\alpha\right)\right)^{\frac{1}{1-q}}$$

The singularity in the normalization factor at $q \to 2$ indicates a phase transition occurring at these parameters.

To express the Lagrange multipliers in terms of the number of particles and energy, we must use the normalization relations (15), in which, in the continuous approximation, we must go from summation to integration and use the density of states proportional $\sqrt{\varepsilon}$ in two qualitatively different cases to $q < 1$ and $q > 1$.

In this case $q < 1$, the nonequilibrium distribution function exists in the interval $0 < \frac{\varepsilon}{T_q} < \frac{1}{1-q}$ and the normalization conditions have the form:

$$\int_0^{\varepsilon_{\max}} d\varepsilon \varepsilon^{1/2} A_q exp_q\left(-\frac{\varepsilon}{T_q}\right) = \eta, \; \int_0^{\varepsilon_{\max}} d\varepsilon \varepsilon^{3/2} A_q exp_q\left(-\frac{\varepsilon}{T_q}\right) = W, \quad (19)$$

For the case $q > 1$, the region of existence of a nonequilibrium distribution function can occupy the entire region of positive energies, and the normalization conditions have the form:

$$\int_0^{\infty} d\varepsilon \varepsilon^{1/2} A_q exp_q\left(-\frac{\varepsilon}{T_q}\right) = \eta, \; \int_0^{\infty} d\varepsilon \varepsilon^{3/2} A_q exp_q\left(-\frac{\varepsilon}{T_q}\right) = W \quad (20)$$

When calculating the resulting integrals, it is convenient to use the following integral relation:

$$I_k(q, w) = \int_0^{w} dz\, z^{k+1/2} exp_q(-z) = \frac{w^{k+3/2}}{k+3/2}\, {}_2F_1\left(k + 3/2, \frac{1}{q-1}, k + 5/2, (1-q)w\right)\Big|$$

The asymptotics of these integrals are also required:

$$b_k(q) = I_k(q, w)|_{w\to\infty} = \frac{\Gamma\left(\frac{3}{2} + k\right)\Gamma\left(\frac{1}{q-1} - \left(\frac{3}{2} + k\right)\right)}{(q-1)^{\frac{3}{2}+k}\Gamma\left(\frac{1}{q-1}\right)}$$

Using these integrals, we write the normalization relations in the form:

$$A_q T_q^{3/2} I_0\left(q, \frac{\varepsilon_{\max}}{T_q}\right) = \eta, \; A_q T_q^{5/2} I_1\left(q, \frac{\varepsilon_{\max}}{T_q}\right) = W$$

Using these notations, we obtain the normalization conditions:

$$T_q = k_{T_q}(q)\frac{W}{\eta},\ A_q = k_{A_q}(q)\frac{\eta}{\left(\frac{W}{\eta}\right)^{3/2}},\ w_{\max} = \frac{\varepsilon_{\max}}{T_q} = \begin{cases} w_{\max} = \frac{1}{1-q},\ q < 1 \\ w_{\max} \to \infty,\ q \geq 1 \end{cases}$$

$$k_{T_q}(q) = \frac{I_0(q, w_{\max})}{I_1(q, w_{\max})} \approx 2.33 - 1.66\,q$$

$$k_{A_q}(q) = \frac{(I_1(q, w_{\max}))^{3/2}}{(I_0(q, w_{\max}))^{5/2}} \approx \exp\left(\frac{2.796}{(1.4-q)^{0.216}} - 2.621\right) \tag{21}$$

The dependence $k_{T_q}(q)$ is close to linear, and $k_{A_Q}(q)$ has a singularity at $q = 1.4$ and is shown in Figure 3.

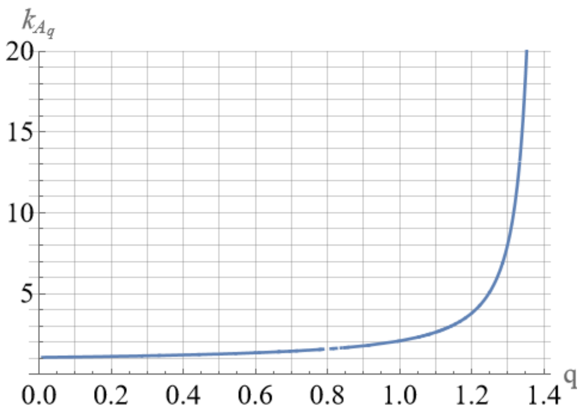

Figure 3. Dependence $k_{A_Q}(q)$ in the expression for the number of particles. When $q = 1.4$ there is a singularity.

By substituting the expressions for $T_q(\alpha, \beta)$ and $A_q(\alpha, \beta)$, one can obtain expressions for and for the Lagrange multipliers.

In the direction parallel to acceleration, collective processes proceed with decreasing spatial scales $q_{\|} \leq 1$, energy values $\frac{m\,u^2}{2}$ are discrete, limited from above, but $T_{\|}$ limited from below and can increase:

$$f_{\|}\left(w_{\|}, T_{\|}, q_{\|}\right) = A_{\|}\left(q_{\|}, T\right) exp_{q_{\|}}\left(-\frac{w_{\|}}{T_{\|}}\right). \tag{22}$$

In the direction orthogonal to acceleration, spatial scales sharply increase (tend to infinity) $q_{\perp} \geq 1$, energy $\frac{m\,{u_{\perp}}^2}{2}$ and continuous quantities that can change in any unlimited intervals and $T_{\perp}$ can both increase indefinitely and tend to zero:

$$f_{\perp}(w_{\perp}, T_{\perp}, q_{\perp}) = A_{\perp}(q_{\perp}, T_{\perp}) exp_{q_{\perp}}\left(-\frac{w_{\perp}}{T_{\perp}}\right) \tag{23}$$

Under conditions approaching equilibrium $q \to 1$, the anisotropy disappears. The greater the degree of nonequilibrium of the system (reflected in the deviation of the nonequilibrium parameter q from 1), the more pronounced the properties of coherently correlated states become.

Thereby, the emergence of anisotropy of the non-equilibrium state of shell type due to the growth of flow in the phase space P_S, happens, naturally, as a result of the simultaneous development of two related processes: localisation in the direction of coherent acceleration with parameter $q \approx \sqrt{1 - P_S}$ as well as delocalisation in the orthogonal direction with the parameter $q \approx \sqrt{1+P_S}$. Wherein, naturally, given that the disequilibrium level decreases, happens the transition towards the symmetrical equilibrium state: $q \underset{P_S \to 0}{\to} 1$

4. Induced Acoustic Plasma Oscillations in Semiconductors in Coherently Correlated States and the Superconducting Transition Temperature

The presence of sources of nonequilibrium in the energy region, which is much higher than the average energy in the equilibrium case, can lead to a pronounced two-component distribution function in terms of energy. As a model for a distribution function of this type in our study, we take a two-component distribution function consisting of an equilibrium distribution function in the region of energies of the order of the Fermi energy (for velocities $u \approx u_F$) and a power-law distribution function in the region of high energies.

Due to the fact that the equilibrium distribution function is $f_0(u, E_F, T_0)$ small in the region $u >> u_F$, we will use the model expression as an analytical representation of the distribution function:

$$f(u) = f_0(u, E_F, T_0) + \eta f_s(u, u_s, q),\ f_s(u, u_s, q) = A_q exp_q\left(-\frac{u^2}{u_s^2}\right).$$

In an equilibrium plasma consisting of two or more groups of charged particles, weakly damped acoustic plasma waves can propagate. Acoustic plasma oscillations propagating in thin semiconductor films were studied in [49], and acoustic plasma oscillations in nonequilibrium states of a semiconductor under the action of microwave radiation were obtained in [50]. Using the distribution function model presented above, we analyze the induced plasma oscillations in a semiconductor. In this model, the permittivity is represented as [56,57]: $\varepsilon^{\updownarrow}\left(\omega, \vec{k}\right) = 1 + \mathrm{Re}\delta\varepsilon + i\mathrm{Im}\delta\varepsilon$, где мнимая часть $\mathrm{Im}\delta\varepsilon = -\frac{8\pi^3 e^2}{m\omega^2}\left(u_p\right)^3 f\left(u_p\right)$ is directly expressed in terms of the distribution function and the phase velocity $u_p = \frac{\omega}{k}$, and the real part by the relation:

$$\mathrm{Re}\delta\varepsilon = -\frac{\omega_{ps}^2}{k^2 u_s^2}\frac{4\pi u_s}{n_s}\int_0^\infty dz \frac{z^2 f(z)}{\nu_p^2 - z^2} \tag{24}$$

where $\omega_{ps}^2 = \frac{4\pi n_s e^2}{m_e}$, $\nu_p = \frac{u_p}{u_s}$, $z = \frac{u}{u_s}$, $u_s = \sqrt{\frac{T_q}{T_0}} u_{T_0}$ are average speed of the nonequilibrium part. Contribution to $\mathrm{Re}\delta\varepsilon$ from the equilibrium part $f_0(u, E_F, T_0)$ has a normal look $-\frac{\omega_p^2}{\omega^2}$, and to calculate the contribution of the nonequilibrium part $f_s(u, q)$ we use the integral [58,59]:

$$I_u\left(u_p, u, q\right) = \int du \frac{u^2 f_s(u,q)}{u_p^2 - u^2} = \chi_q\left(u_p, u\right) {}_2F_1\left(1, \frac{2-q}{1-q}, \frac{3-2q}{1-q}, \frac{1+(q-1)u^2}{1+(q-1)u_p^2}\right)$$

where $\chi_q\left(u_p, u\right) = \frac{1}{2}\frac{(q-1)\left(1+(q-1)u^2\right)^{\frac{2-q}{1-q}}}{(2-q)\left(1+(q-1)u_p^2\right)}$. If we use asymptotics (22), we get that $\varepsilon^{\updownarrow}\left(\omega, \vec{k}\right)$ looks like:

$$\varepsilon^l\left(\omega, \vec{k}\right) = 1 + \frac{k^2}{k_D^2} - \frac{\omega_{p0}^2}{\omega^2} + \frac{k_{eff}^2(q)}{k^2},\ k_{eff}(q) = \frac{\omega_{ps}}{u_{eff}(q)}, \tag{25}$$

$u_{eff}(q) = \langle u(q)\rangle$ is velocity averaged over the nonequilibrium distribution function with the nonequilibrium parameter q. Averaging using the integral $I_k(q, w)$ leads to the dependence on the nonequilibrium parameter shown in Figure 4.

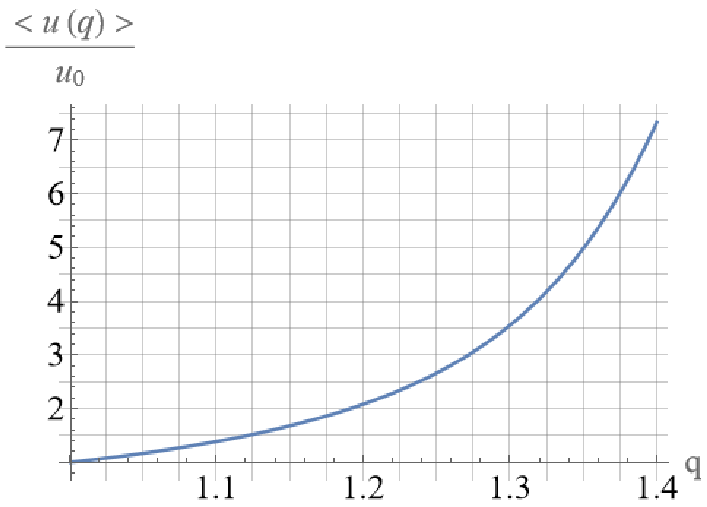

Figure 4. Dependence of the average velocity $u_{eff}(q) = \langle u(q) \rangle$ on the nonequilibrium parameter q.

From (25) and $\varepsilon^l\left(\omega, \vec{k}\right) = 0$ follows the dispersion law of longitudinal vibrations:

$$\omega^2 = \omega_p^2\left(1 + \frac{k^2}{k_D^2} + \frac{k_{eff}^2(q)}{k^2}\right)^{-1}.$$

Dispersion of longitudinal vibrations in the jelly model for different nonequilibrium parameters is shown in Figure 5.

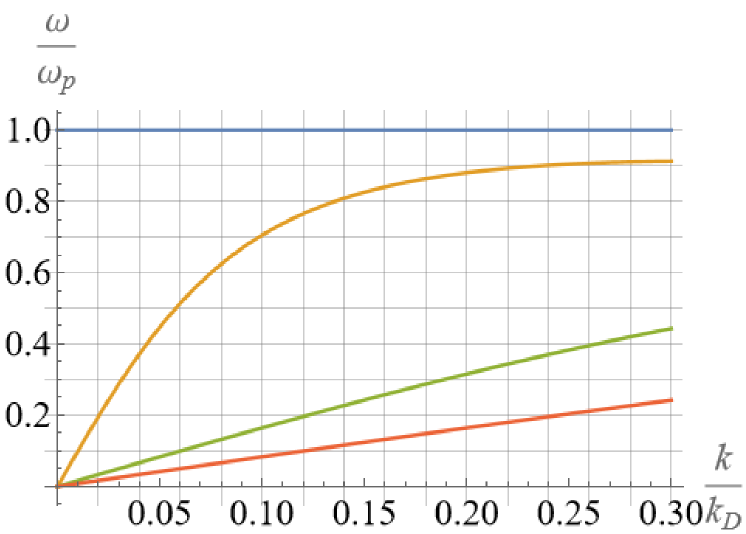

Figure 5. Dispersion of longitudinal vibrations in the jelly model for different nonequilibrium parameters. Above is a straight line corresponding to the equilibrium case, and below are curves for the values $\frac{k_{eff}(q)}{k_D} = 0.1,\ 0.6,\ 1.2$, respectively.

The influence of the medium on the interaction between electrons and, accordingly, on the phase transition temperature is conveniently represented by the functional of the permittivity [25], a set of branches of longitudinal vibrations $\Omega_k = \left(\omega_j(k), E_F\right)$ and force constants:

$$a_j = \left(\omega^2 \frac{\partial \varepsilon}{\partial \omega^2}\right)^{-1}_{\omega^2 = \omega_j^2(k)},\ \overline{\xi} = \prod_j \omega_j^{\left(a_j\left(\sum_j a_j\right)^{-1}\right)}$$

$$T_c \approx \prod_k \Omega_k^{A_k \left(\sum\limits_k A_k\right)^{-1}} \exp\left(-\left(\mu \sum_k A_k\right)^{-1}\right), \tag{26}$$

when there is only one branch of longitudinal oscillations, then for the force constant, it follows:

$$a = \frac{\omega^2(\bar{k})}{\omega_p^2}\left(1 + \frac{\omega_{ps}^2}{\bar{k}^2 u_{eff}^2}\right)^{-1} \approx \left(\frac{\bar{k}^2 u_{eff}^2(q)}{\omega_{ps}^2}\right)^2,$$

where $\bar{k} = \alpha k_F$, $u_{eff} = \beta u_F$ and, accordingly, for T_c:

$$T_c = \hbar\omega_p \exp\left(-\frac{1}{\mu(a-1)}\right),\ a = \left(\frac{n}{n_s}\right)\left(E_F/\hbar\omega_p\right)^2 \alpha^2 \beta^2 \tag{27}$$

$$\frac{T_c}{\hbar\omega_p} = \exp\left(-\left(\mu\left(\left(\frac{n}{n_s}\right)\left(\frac{E_F}{\hbar\omega_p}\frac{u_{eff}(q)}{u_F}\frac{\bar{k}}{k_F}\right)^2 - 1\right)\right)^{-1}\right)$$

Figure 6 shows the dependence $\frac{T_c}{\hbar\omega_p}$ on the nonequilibrium parameter q.

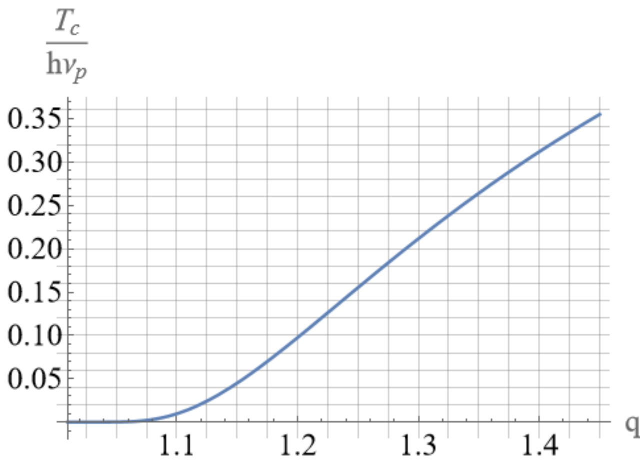

Figure 6. Dependence $\frac{T_c}{\hbar\omega_p}$ on the nonequilibrium parameter.

Thus, both from the analytical expression for the superconducting transition temperature and from Figure 6, it can be seen that with increasing nonequilibrium, the phase transition temperature increases strongly and tends to the energy of the quantum of electronic plasma oscillations.

5. Conclusions

The paper presents the Vlasov kinetic equations and the corresponding dynamic transfer equations after averaging over higher accelerations, considering the change in the internal structure and, consequently, the inertia of the system. The obtained exact solutions of the Vlasov equation, which take into account the variation in the structure of the system, are given in the form of non-extensive Tsallis distribution functions and are consistent with the exact solutions of the kinetic equations describing quasi-homogeneous systems with flows in the phase space, and the thermodynamic distributions, which are obtained together with the corresponding expressions for the Lagrange coefficients—nonequilibrium temperature and number of states depending on the order parameter and nonequilibrium coefficient for coherently correlated states.

For the model distribution function, which reflects the order parameter in the system and its strong non-equilibrium, the phase transition temperature in the "jelly" model was estimated using V. Ginzburg's functional relations. Using the thermodynamic relations obtained in the work, the dependence of the transition temperature T_c on nonequilibrium parameters q is given, and it is shown that with an increase in the nonequilibrium parameter q, the phase transition temperature T_c increases strongly and tends to the energy of the of electron plasma oscillations quantum $T_c \rightarrow \hbar\omega_p$ under increasing $q > 1$ at the nonequilibrium coherently correlated states.

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