Higher order long range correlations in nuclear structure and dynamics

Vladimir N. Kondratyev [1,2,*] and Massimo Di Toro [1]
[1] *Laboratorio Nazionale del Sud, INFN, v.S.Sofia 44, 95125 Catania, Italy*
[2] *Fakultät für Physik, Bielefeld Universität, D-33501 Bielefeld Germany*
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An explicit correction term to the ensemble averaged trajectory for the phase space single-particle distribution function due to the presence of the higher order long range correlations is considered. It is demonstrated that this extension of the usually employed transport approaches (BUU, BNV, LV, etc.) has a diffusion structure. The role of higher order correlations in the cases of nuclear collective motion and the mean field decomposition effect in nuclear fragmentation at high temperatures is analyzed. In particular we show the importance of higher order correlations for the transition from zero to first sound regime for propagation of the collective excitation in hot nuclear matter. [S0556-2813(96)00905-3]

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In the last decade the use of the kinetic transport theories of the Boltzmann-Nordheim-Vlasov (BNV) type has been very successful in describing the main reaction mechanisms (e.g., particle production, collective flows, etc.) of heavy ion collisions (for review see, for example, [1]) All these theories are based on the one-body reduction of the many-body dynamics accounting for the two-body correlation effects through the averaged response of the one-body density [long range correlations (LRC)] and the binary collision term [short range correlations (SRC)]. At sufficiently high excitation energies the dynamics of many-body correlations can bring new important information on the system, going from the variances of physical observables to the essential features of the time evolution in instability regions [1,2]. Recently there have been proposed several attempts to formulate a stochastic transport theory including a selfconsistent dynamics of the higher order correlations. These approaches are based on the Boltzmann-Langevin (BL) method, where the higher order effects are treated through the fluctuations associated with the collision integral (e.g., the stochastic part of the SRC [3–5]). In some previous studies regarding the intermediate excitation energies (see, for example, [6–9], and references therein) it has been pointed out the importance of higher order long range correlations (HOLRC), that can be treated as a mean field fluctuations. In this paper such effect is analyzed within a technique which is not restricted to the usually employed mode expansion approximation. We show that being more general such a method gives some simplifications in understanding the physics and gross properties of HOLRC effects. In particular we discuss here an application to cases where we expect an increase of the contributions from long range correlations and a relatively large spreading of different events: the collective motions at high temperature and the dynamical instabilities. We must remark that many studies have been devoted to the analysis of higher order correlation effects in the theory of nuclear collective motions, in particular of giant resonances [10–13]. It is however extremely difficult to predict the temperature dependence of these contributions. In this paper we suggest a semiclassical approach which allows one to get a quite clear picture of correlation contributions at high temperature with fundamental consequences on the nature of collective motions in excited nuclear matter.

The general features of the kinetic phenomena can be expressed in terms of the average properties of the single-particle (sp) distribution function (df): $f(t)=\langle\hat{f}(t)\rangle$; where $\langle\cdots\rangle$ refers to the ensemble averaging related, for example, to nucleus-nucleus collision events. The equation of motion for the fluctuating sp df $\hat{f}(t)$ regarding the particular event can be written in a very general way as

$$\frac{\partial\hat{f}}{\partial t}+iL\hat{f}=I(\hat{f})+\delta K(t), \tag{1}$$

where $L=-i\{h,\ldots\}$ is the Liouville operator corresponding to an evolution of the system in ensemble averaged time dependent mean field $U(\rho)=\langle\hat{U}(\hat{\rho})\rangle$, and $h(\rho)=p^2/2m+U(\rho)$ is the self-consistent mean field Hamiltonian, related to the long range part of the nuclear interaction. $I(\hat{f})$ denotes, in general, a memory dependent binary collision term, representing the average effect of the residual Pauli-reduced hard two-body interactions. This gives rise to the volume dissipation due to incoherent two-body correlation effects during the evolution. The higher order effects are indicated in Eq. (1) by the additional term: $\delta K(t)=\delta K_{\rm col}(t)+\delta K_{\rm mf}(t)$; corresponding to the fluctuating part of the SRC: $\delta K_{\rm col}(t)$; as well as to the HOLRC: $\delta K_{\rm mf}(t)=-i\,\delta L\hat{f}$; $\delta L=-i\{\delta U,\ldots\}$, where $\delta U=\hat{U}-U$ is the deviation of the mean field from the average one. This δU has nothing to do with variations of the average mean field due to self-consistency. These are indeed already accounted for within standard random-phase approximation (RPA) approach to the linear response theory for the ensemble averaged sp df (see below and in the Appendix). Moreover we show below that contrary to the versions of BL model incorporated in [4, 9] such a fluctuation is not vanishing in a collisionless dynamics and in a zero sound propagation of collective motions. Actually one of the main results of the present study is

*Permanent address: Institute for Nuclear Research, 47, Pr.Nauki, Kiev, 252028 Ukraine.

related to the effect of the HOLRC on the transition from zero to first sound dynamical regime.

We consider the fluctuating properties of the nuclear dynamics in analogy with Brownian motion, where it is assumed that Eq. (1) describes a stochastical process in which the entire sp df $\hat{f}$ is a stochastic variable and $\delta K(t)$ acts like a random force. For the intermediate excitation energies (above 20 MeV) it is convenient to employ the Markovian assumption. Then the closed set of the self-consistent transport equations for $f(t)$ is derived (see the Appendix) from Eq. (1) to be

$$\frac{\partial f}{\partial t}+iM[f]f=K_c[f], \quad M[f]\equiv L[\rho]+iI[f], \tag{2}$$

$$K_c[f]=\sum_{\alpha\beta} D_p^{\alpha\beta}\partial_{p\alpha}\partial_{p\beta}f(\mathbf{r},\mathbf{p};t), \tag{3}$$

$$\frac{d}{dt}\sigma_{\rm var}=-i(M[f]-M^*[f])\sigma_{\rm var}+D(t), \tag{4}$$

$$D(t)\equiv\sum_{\alpha\beta} D_p^{\alpha\beta}\partial_{p\alpha}f\ \partial_{p\beta}f. \tag{5}$$

Projecting Eq. (4) to some particular modes one gets [14] the so-called Lalime equation for the expansion of the correlation function over the collective vibration modes.

The diffusion term (3) and (5) can be further simplified in the case near local equilibrium. This case is of interest for studies of higher order correlation effects in the nuclear collective dynamics (e.g., giant resonances, fusion, fission, and fragmentation phenomena). In this quasistatic limit with respect to the diffusion properties Eq. (4) reduces to

$$D_{\rm eq}=i(M-M^*)_{\rm eq}\sigma_{\rm var}^{\rm eq}=2\sigma_{\rm var}^{\rm eq}/\tau_{\rm rel}, \tag{6}$$

where $\tau_{\rm rel}$ is a relaxation time for nucleon dynamics representing the composition of the collisional and sp mechanisms: $\tau_{\rm rel}^{-1}=\tau_{\rm col}^{-1}+\tau_{\rm sp}^{-1}$, accounted by an extended sp evolution operator M [15]. One sees that an additional correlation term due to the coupling sp dynamics with mean field fluctuations is not vanishing also within the limit of the collisionless regime $\tau_{\rm col}^{-1}\to 0$. Using the familiar result $\sigma_{\rm var}^{\rm eq}=f_{\rm eq}(1-f_{\rm eq})\approx T df/dh$; obtained within semiclassical mean field approximation neglecting zero point vibrations, and employing the definition (5) we get from Eq. (6) the famous Einstein relation for the average diffusion rate: $\overline{D_p^{\rm eq}}\approx 2T/\overline{\tau_{\rm rel}}$.

One of the particularly important quantities related to the dynamical and structure features is represented by the sp linear response function (rf). Using the perturbation theory (pt) and standard semiclassical technique (see, for example, [16]) we obtain from Eqs. (2) and (3) the following expression for the imaginary part of the sp rf: $\chi_0''(\mathbf{k};\omega)\equiv{\rm Im}(\chi_0(\mathbf{k};\omega))$, of the Fermi system

$$\chi_0''(\mathbf{k};\omega)=\frac{\omega}{\pi}\int dt\ C(\mu,t)\exp(-i\omega t), \tag{7}$$

with

$$C=\int\frac{d\mathbf{r_1}d\mathbf{p}}{(2\pi)^3}\delta(\mu-h)\Sigma(\mathbf{r_1},\mathbf{r};t)\exp\{i\mathbf{k}(\mathbf{r}-\mathbf{r}_1)\}, \tag{8}$$

where μ denotes the nucleon chemical potential

$$\mu\approx\epsilon_F\left(1-\frac{5\pi}{4}\frac{\epsilon^*}{\epsilon_F}\right)$$

with $\epsilon^*\approx\pi^2T^2/4\epsilon_F$; $\Sigma(\mathbf{r_1},\mathbf{r};t)$ is the probability of finding a nucleon with initial conditions $(\mathbf{r},\mathbf{p})$ at the point $\mathbf{r_1}$ at the time t. Assuming $T\ll\epsilon_F$, we have employed here a kind of zero-temperature approximation that consists in the following relation: $df/dh=(2\pi)^{-3}2\delta(\mu-h)$. It is important to stress here that the approximation (7) is valid in the limit of the long wavelength field $k\ll k_\mu=\sqrt{2m\mu}$.

Within the local density approximation the diffusion rate has a scalar form $D_p^{\alpha\beta}\approx\delta_{\alpha\beta}D_p$. Then the probability $\Sigma(\mathbf{r_1},\mathbf{r};t)$ can be approximated by the Gaussian distribution function (see [19]):

$$\Sigma(\mathbf{r_1},\mathbf{r};t)\approx\frac{\exp\{-[\mathbf{r_1}-\mathbf{r}(t)]^2/D_rt\}}{(\pi D_rt)^{3/2}}, \tag{9}$$

where $\mathbf{r}(t)$ denotes the average trajectory of a nucleon and the spatial diffusion coefficient D_r is related to the value $\overline{D_p}$: $D_r=\overline{\tau_{\rm rel}^2}\overline{D_p}$. For infinite nuclear matter we assume rectilinear average trajectory $\mathbf{r}(t)=\mathbf{r}+\mathbf{v}t$; and calculate from Eqs. (7)–(9) the quantity $\chi_0''(\mathbf{k};\omega)$ as [16,20]

$$\chi_0''(z,u)\approx\frac{mp_\mu}{2\pi^2}u\Bigg(\Theta(1-(z\Delta)^2-u^2)+\pi^{-1}\arctan\left(\frac{2z\Delta}{(z\Delta)^2-1+u^2}\right)\Bigg), \tag{10}$$

where the variables z and u are related to the wave number $(z=k/2k_\mu)$ and sound velocity $(u=\omega/kv_\mu)$ measured in the respective units regarding the chemical potential, while $\Delta=D_r/2\hbar$.

The real part of the sp rf $(\chi_0'(\mathbf{k};\omega))$, is calculated using the Kramers-Kronig relation

$$\chi_0'(z,u)=\frac{mp_\mu}{2\pi^3}(1-\frac{u}{4}\ln\left(\frac{(u+1)^2+(z\Delta)^2}{(u-1)^2+(z\Delta)^2}\right)). \tag{11}$$

An expansion of the response function for small values of product $(z\Delta)$, which represents the ratio of the in medium correlation length and the wavelength, reads

$$\chi_0(z,u)|_{(z\Delta)\ll 1}\approx\frac{mp_\mu}{2\pi^3}\left(1-\frac{u}{2}\ln\left|\frac{u+1}{u-1}\right|+i\pi u\Theta(1-u)+i\frac{2uz\Delta}{u^2-1}+\frac{(uz\Delta)^2}{(u^2-1)^2}+\cdots\right). \tag{12}$$

From first three terms of this series one can easily recognize the familiar result obtained within Vlasov approximation [17,18]. We see that HOLRC are extremely important if the reduced sound velocity is close to 1. It is also clear that the diffusion term acts in the direction of a reduction of the real

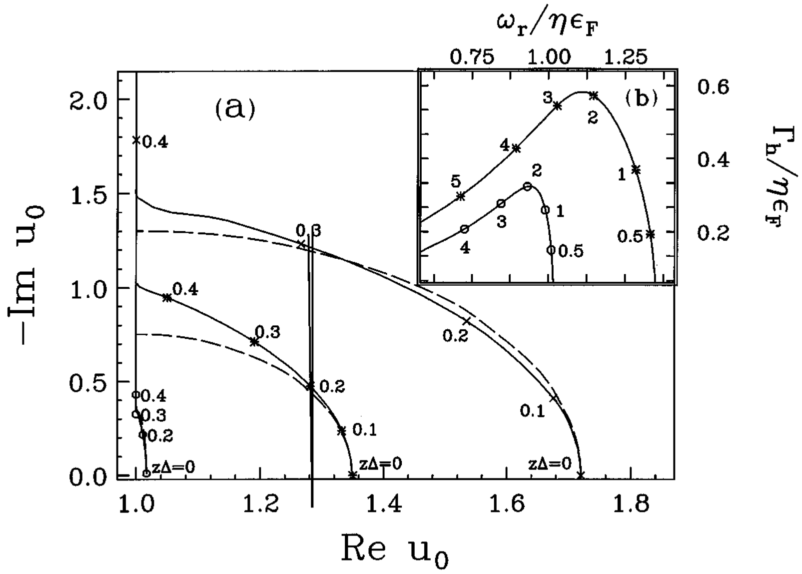

FIG. 1. (a) The sound velocity u_0 vs the product $(z\Delta)$: the solid line is the solution of Eq. (13); the dashed line is obtained according to Eq. (14) with $\kappa_i \approx 1.085$, 2.38, and 4.71 for $F_0 = 0.7$ (○); 3.5 (*) and 7.0 (×), respectively. (b) The temperature dependence of the reduced GDR frequency and HOLRC partial width being the solution of Eqs. (13) and (15) at $\zeta = 25$ and $\kappa_0 = \omega_r^2/\Gamma_{\rm GDR}^{\rm col} \approx 225$ MeV, corresponding to mass number $A \approx 120$. The figures on the curves are in units: $(10\ T)/\eta\epsilon_F$.

part of the sound velocity. Such a situation occurs, for instance, in the case of zero sound (see below).

The sp response function allows one to consider the characteristic vibrations representing important features related to the structure of the many-body system. These vibrations are determined by the following dispersion relation [17,18]:

$$\chi_0(\mathbf{k},\omega_{\mathbf{k}}) = -\left(\frac{\partial U_{\mathbf{k}}}{\partial \rho}\right)^{-1}. \tag{13}$$

Here the coupling constant is given by $\partial U_{\mathbf{k}}/\partial\rho = g(\mathbf{k})(\partial\widetilde{U}/\partial\rho)$, where $\widetilde{U}$ is a function of the local density ρ and $g(\mathbf{k})$ is the Fourier component of the convolution function that simulates a finite range of the interaction.

A very important property of a Fermi liquid is related to the possibility of propagation of zero sound waves at very low temperature [17]. These vibrations are associated with a solution of Eq. (13) with real value of sound velocity corresponding to the condition $\mathrm{Re}[u] > 1$. It is well known that in this case we have $(\omega\tau_{\rm col}) > 1$ and the dominant relaxation mechanism at low temperatures turns out to be the two-body one due to the collision term [see Eq. (2)]. In this case the relaxation rate regarding the collisional width of the collective mode grows with the temperature as T^2. It can be easily seen that the solution associated with zero sound can exist if the value $F_{\mathbf{k}} = (\partial U_{\mathbf{k}}/\partial\rho)(3\rho/2\epsilon_F)$ is positive. This situation occurs for nuclear matter at density values larger or equal to the normal one. For our estimates we consider a wave number independent parameter F_0. From Fig. 1 we see that an additional diffusion term in the transport equation due to the HOLRC results in the decrease of the real part of sound velocity. On the other hand, the imaginary part grows proportionally to the product $z\Delta$. This behavior can be approximated by

$$\frac{\mathrm{Re}[u_0]-1}{\overline{u}_0-1} \approx \sqrt{1-\left(\frac{z\Delta}{\kappa_r}\right)^2}, \quad \mathrm{Im}\,[u_0] \approx \kappa_i(z\Delta), \tag{14}$$

where $\overline{u}_0$ is the sound velocity related to the respective Vlasov equation, $\kappa_r \approx 0.32$; and $\kappa_i \approx 1 + 2\exp\{-2/F_0\}$ for small values of F_0. The imaginary part of sound velocity is associated with a damping of the collective mode due to correlations. Therefore in the case of comparable correlation and wavelengths, $(z\Delta) \sim 1$, we can come to a situation when the HOLRC contribution to the width of the collective vibration is larger than the respective frequency.

For the giant dipole resonance (GDR) in finite nuclei we can estimate the wave number according to Steinwedel and Jensen model: $k_{\rm GDR} \approx \pi/2R$, where R is the nuclear radius. In the bulk region we assume the collisional relaxation mechanism for the nucleons participating to dipole collective mode: $\tau_{\rm rel} \approx \tau_{\rm col} \approx \kappa_0/(\omega_r^2 + \zeta T^2)$, that results in a frequency dependent diffusion rate [see Eq. (6)]. Then the coupled equations accounting for HOLRC are reduced to

$$(z\Delta) \approx \frac{\eta}{2}\,\frac{\kappa_0 T}{\omega_r^2 + \zeta T^2}, \quad \omega_r - i\Gamma_h \approx \eta\mu u_0, \tag{15}$$

where Γ_h indicates the contribution of HOLRC to the resonance width; and $\eta = \pi/k_\mu R \approx 2.1 A^{-1/3}$. We see that the product $(z\Delta)$ reaches its maximum value at temperature $T_{\rm tr} = \omega_r/\sqrt{\zeta}$. Such a behavior could be related to the recently reported experimental evidence dealing with saturation of the GDR width at high excitation energy [21] starting from the temperature $T_{\rm tr} \approx 3$ MeV. This evidence allows one, indeed, to estimate the parameter $\zeta \approx 25$. This value is between the various theoretical evaluations given in the literature: $\zeta = 4\pi^2/3$ and $\zeta = 4\pi^2$ (see [21] and references therein). Figure 1(b) shows that the HOLRC partial width sharply increases with temperature, reaches its maximum value at $T_{\rm tr} \approx \eta\epsilon_F/\sqrt{\zeta}$, and slowly decreases at $T > T_{\rm tr}$. We stress that without the important corrections due to HOLRC the transition temperature would be around 15 MeV [22,21], much more than an equilibrated compound nucleus can sustain. Therefore without the substantial contribution of HOLRC it would be not possible to observe such an effect in excited nuclei. Similar behavior of the sound waves damping has been observed, for example, in ^{3}He liquid [23]. It would be very interesting to carry out similar studies in the case of nuclear matter.

Finally we briefly consider the conditions referring the unstable modes. In this case one finds the imaginary solution $\omega = i\nu$ of the dispersion relation (13) and it is defined by the RHS of the Eq. (10). The negative quantity ν corresponds to a positive Lyapunov exponent of the respective collective mode leading to a mean field (spinodal) decomposition of the nuclear matter with growing time $\tau_k = -(\nu k)^{-1}$. From Eqs. (10) and (13) we find that a solution referring to amplified modes exists at the condition $(1+F_k^{-1})<0$, while the case $(1+F_k^{-1})>0$ is associated with the overdamped modes. Similar conditions have been obtained in [14] also in the limit of Vlasov approximation.

We have presented some results on higher order long range correlation effects in nuclear dynamics. The obtained equations (3)–(5) for the HOLRC term can be used in order to test the mode expansion approximation. Following a semiclassical treatment we have analyzed some gross properties of such an effect considering the averaged behavior of the respective kinetic terms and employing a Markovian process assumption. We have seen that HOLRC's are very important for a correct description of thermal properties of zero sound propagation.

On the other hand, one can find some analogy between the expansion of the response function with respect to the quantity z and quantum features related to the intrinsic dynamics [24]. In absence of HOLRC's, quantum corrections associated with finiteness of the wave-length start with a term proportional to z^2, see Eq. (10) of [24]. From Eq. (12) we see that a term linear in z is not vanishing as well due to the presence of HOLRC's. This property seems to indicate a quantum origin of considered correlations which turns out to be in competition with quantum features of intrisic structure. The mean field fluctuations result in a broadening of single-particle energies and eventually in a melting of the intrinsic quantum structure at high temperature. For sound waves a naive estimate of such ''melting point'' can be obtained from the condition $|u\Delta|>z$. The use of the respective parameters in finite nuclei (see above) yields to a relatively small value of the associated temperature $T>T_m \approx \overline{\tau}_{\rm rel}^{-1}/|u|k_\mu R$. However our results also show that the considered effects require a more detailed analysis when quantum features, like retardation effects, etc., are taken into account as well. A quantum Langevin treatment, similar to the one presented in [9], would allow one, indeed, to carry out such a study, that we hope to present in a near future.

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APPENDIX: THE STOCHASTIC HIGHER ORDER LONG RANGE CORRELATIONS

The transport equation (2) for the averaged df is obtained by statistical averaging of the respective equation (1). Then the correlation term arising on the RHS of Eq. (2) is written as [25]

$$K_c[f] = \langle \delta K(t) \rangle \approx -i \langle \delta L\, \delta f \rangle, \tag{A1}$$

where $\delta f = \hat{f} - f$ is the deviation of the df from the average one. We neglect the fluctuating part of the SRC $[\delta K_{\rm col}(t)]$, that vanishes near the equilibrium due to the averaging. This part could turn out to be important, indeed, in some specific cases dealing with many-body coherent effects within the local region, for example subthreshold particle production.

For the small amplitude fluctuations (δf) we use the pt by employing the linearized equation (1) around the averaged equation (2):

$$\frac{\partial \delta f}{\partial t} + iM[f]\delta f = \delta K(t), \tag{A2}$$

where $M[f]$ represents an extended sp evolution operator including the combined action of the averaged self-consistent mean field and collision integral [see Eq. (2)]. Solving formally Eq. (17) and assuming a random force structure of the fluctuation term,

$$\delta K(t) \approx \sum_\alpha \delta F_\alpha \partial_{p\alpha} f \quad \text{with} \quad \delta F_\alpha \approx -\partial_{r\alpha} \delta U,$$

we get the stochastical correlation term $K_c[f]$ of Eq. (16) as

$$K_c[f] = \int_0^t ds \sum_{\alpha\beta} \left\langle \partial_{r\alpha} \delta U(t) \exp\left\{ -i \int_s^t d\tau\, M^{cl}(\tau) \right\} \partial_{r\beta} \delta U(s) \right\rangle \partial_{p\alpha}\partial_{p\beta} f(s) + \left\langle \vec{\partial}_r \delta U(t) \exp\left\{ -i \int_0^t d\tau\, M^{cl}(\tau) \right\} \vec{\partial}_p \delta f(0) \right\rangle. \tag{A3}$$

Assuming Markovian structure of the correlations and a local mean field we obtain [25] the form (3) for the term $K_c[f]$ with

$$D_p^{\alpha\beta} = \int_0^t ds\, \partial_{r\alpha} \partial_{r_1\beta} \langle \delta U(\mathbf{r},t)\, \delta U(\mathbf{r}_1,s) \rangle |_{\mathbf{r}_1 \to \mathbf{r}}. \tag{A4}$$

One sees that the term accounting for HOLRC has a diffusion structure with diffusion rate related [25] to the density correlation function within to a coupling constant factor [see Eq. (13)]. The dynamics of correlation function is, consequently, affected by HOLRC effects. These properties allow one, indeed, to incorporate the HOLRC in a self-consistent manner.

Employing Eq. (17) for the sufficiently small time interval $[\Delta t = t - t_0]$ we get the df correlation function $\sigma_\Gamma(t_1,t_2) \equiv \langle \delta f(t_1) \delta f^*(t_2) \rangle$ as

$$\sigma_\Gamma(t_1,t_2)=\int_{t0}^{t1}ds\ \exp\left\{-i\int_s^{t1}d\tau\ M\right\}\int_{t0}^{t2}ds'\exp\left\{-i\int_{s'}^{t2}d\tau\ M*\right\}\langle\delta K(s)\,\delta K^*(s')\rangle$$
$$+\exp\left\{-i\int_{t0}^{t1}d\tau\ M\right\}\exp\left\{-i\int_{t0}^{t2}d\tau\ M\right\}\sigma_\Gamma(t_0,t_0). \quad \text{(A5)}$$

Under the Markovian approximation:

$$\langle\delta K(s)\,\delta K(s')\rangle\approx\sum_{\alpha\beta}\langle\partial_{r\alpha}\delta U(s')\,\partial_{r\beta}\delta U(s)\rangle\partial_{p\alpha}f\,\partial_{p\beta}f\approx\sum_{\alpha\beta}D_p^{\alpha\beta}\partial_{p\alpha}f\,\partial_{p\beta}f\,\delta(s-s')\equiv D(s)\,\delta(s-s') \quad \text{(A6)}$$

the equal time correlator $\sigma_{\rm var}(t)\equiv\sigma_\Gamma(t,t)$ turns out to be dominant and it has the following form:

$$\sigma_{\rm var}(t)=\exp\left\{-i\int_{t0}^{t}d\tau(M-M^*)\right\}\sigma_{\rm var}(t_0)+\int_{t0}^{t}ds\ \exp\left\{-i\int_s^t d\tau(M-M^*)\right\}D(s). \quad \text{(A7)}$$

One can easily see that the phase space correlation function fulfills the differential equation (4).

[1] G.F. Bertsch and S. Das Gupta, Phys. Rep. **160**, 190 (1988); A. Bonasera, F. Gulminelli, and J. Molitoris, *ibid.* **243**, 1 (1994).
[2] U. Mosel, Nucl. Phys. **A583**, 29c (1995).
[3] S. Ayik and C. Gregoire, Phys. Lett. B **212**, 269 (1988); Nucl. Phys. **A513**, 187 (1990).
[4] J. Randrup and B. Remaud, Nucl. Phys. **A514**, 339 (1990).
[5] F. Chapelle, G.F. Burgio, Ph. Chomaz, and J. Randrup, Nucl. Phys. **A540**, 227 (1992), and references therein.
[6] Y. Alhassid and B. Bush, Phys. Rev. Lett. **63**, 2452 (1989); Nucl. Phys. **A514**, 434 (1990); **A531**, 1 (1991).
[7] F.V. De Blasio, W. Cassing, M. Tohyama, P.F. Bortignon, and R.A. Broglia, Phys. Rev. Lett. **68** 1663 (1992).
[8] A. Bonasera, M. Colonna, M. Di Toro, F. Gulminelli, and A. Smerzi, Nucl. Phys. **A572**, 171 (1994).
[9] S. Ayik, Z. Phys. A **350**, 45 (1994).
[10] J. Speth and J. Wambach, International Review of Nuclear Physics (World Scientific, Singapore, 1991), Vol. 7, p. 1.
[11] G.F. Bertsch, P.F. Bortignon, and R.A. Broglia, Rev. Mod. Phys. **55**, 287 (1989).
[12] V.G. Soloviev, Z. Phys. A **334**, 143 (1989); Prog. Part. Nucl. Sci. **28**, 49 (1992), and references therein.
[13] V.V. Voronov and V.Yu. Ponomarev, in *Selected Topics in Nuclear Structure*, edited by V.G. Soloviev (JINR, Dubna, 1994), p. 9.
[14] M. Colonna, Ph. Chomaz, and J. Randrup, Nucl. Phys. **A567**, 637 (1994); M. Colonna and Ph. Chomaz, Phys. Rev. C **49**, 1908 (1994).
[15] G.F. Burgio and M. DiToro, Nucl. Phys. **A476**, 189 (1988); V.M. Kolomietz and V.N. Kondratyev, Z. Phys. A **344**, 125 (1992).
[16] V.N. Kondratyev, Phys. Lett. A **179**, 209 (1993); **190**, 465 (1994).
[17] L.D. Landau, Sov. Phys. JETP **5**, 101 (1957).
[18] I.M. Khalatnikov and A.A. Abrikosov, Sov. Phys. JETP **6**, 84 (1958).
[19] K.M. Rattray and A.J. McKane, J. Phys. A **24**, 4375 (1991); G. Munoz, J. Phys. A **26**, 4475 (1993).
[20] V.N. Kondratyev and A. Bonasera, Phys. Rev. Lett. **74**, 2824 (1995); V.N. Kondratyev and V.A. Shchepunov, Radiat. Meas. **25**, 105 (1995).
[21] M. Di Toro *et al.*, Proceedings International Conference on Giant Resonances, Groningen, 1995 [Nucl. Phys. **A** (to be published)], and references therein.
[22] M. Colonna, M. DiToro, V. Latora, and A. Smerzi, Prog. Part. Nucl. Sci. **30**, 17 (1993).
[23] W.R. Abel, A.C. Anderson, and J.C. Wheatley, Phys. Rev. Lett. **17**, 74 (1966).
[24] S. Ayik, M. Colonna, and Ph. Chomaz, Phys. Lett. B **353**, 417 (1995).
[25] V.N. Kondratyev, Z. Phys. B **99**, 473 (1996).