Contribution of density fluctuations to the damping of giant resonances

K. Morawetz
Max-Planck-Gesellschaft, AG ''Theoretical Many-Particle Physics,'' University Rostock, 18051 Rostock, Germany

M. Di Toro
Laboratorio Nazionale del Sud, Viale a. Doria - 95123 Catania, Italy
(Received 23 January 1996)

A generalized kinetic equation of Lennard-Balescu type is derived with retardation effects. By explicit consideration of the density-density fluctuations we give the contribution of fluctuations to the spreading width of giant resonances besides the collisional part. For small frequencies and long wavelengths, approximate formulas for the damping of giant resonances are derived. The density fluctuations couple the equation of state to the damping rate and give rise to an overall enhancement factor. This density and temperature dependent factor can reproduce the experimental values of the damping. Inside the spinodal region the factor turns out to be negative, indicating unstable modes. Explicit comparisons are made with the experimental values of giant monopole and dipole resonances. It is shown that the experimental widths can be described assuming finite nuclei since surface effects are important for the influence on density fluctuations. The recently reported saturation of the dipole width with increasing energy is explained as a transition between zero and first sound. Within this treatment we find an overcritical range of fluctuations, above which the nucleus is destroyed and no collective behavior is possible at all. [S0556-2813(96)02307-2]

PACS number(s): 24.30.Cz, 24.60.Ky

I. INTRODUCTION

Since the first discovery of collective vibrations in nuclei by Baldwin and Klaiber in 1947 [1] the giant resonances have been studied intensively. Most data are available for giant dipole resonances [2,3], which corresponds to a vibration out of phase of protons against neutrons. It is found that this resonance is a property of all nuclei. The restoring force is generated by the attractive nuclear force between the nucleons. Thus it is an important experimental signal to study the microscopic nuclear forces. Restoring forces currently provide quite reliable information on bulk properties like compression moduli, symmetry energy, spin-isospin sound velocity, etc., of symmetric and asymmetric nuclear matter. Recent experimental results concern double giant dipole resonances [4], i.e., resonances built on resonances.

Many experimental measurements were performed by photoabsorbtion reactions [5,6]. There the giant resonance can well be described by a Lorentzian distribution whose centroid energy and width vary smoothly with nuclear mass. This indicates that these are collective phenomena in which many nucleons participate [7,8]. Collective vibrations occur quite universally in interacting systems forming clusters [9], like in liquid metals or nuclei. Therefore it is one of the challenging questions to understand the width of such giant resonances, because this can give information on the energy dissipation in Fermi liquids. A lot of attempts are made to understand these properties within a microscopic picture [10,11] (and citations therein).

From optical potential analysis of nucleon-nucleus scattering at low energies one can deduce that the nucleon mean free path is of the order of the nuclear radius. Thus the quasiarticle picture is an appropriate starting point for the description of giant resonances. Therefore the kinetic description is believed to be applicable. This leads to the picture of quasiparticles moving through a common mean field and being described by the collisionless Vlasov equation. To describe the spreading width one has to account for collisions. This is done by linearization of the Vlasov equation which yields the random phase approximation (RPA) for small amplitudes [12–15]. In this way RPA theories can describe the giant resonances correctly [11]. The solution of these collective modes within the Vlasov approach can be found in [16–18]. In principle, one goes beyond collisionless transport in allowing energy dissipation, that is definitely not the case in the Vlasov equation, which is reversible. Therefore the extension of the Vlasov equation represents already a higher-order type of collision integral, i.e., the Lennard-Balescu collision integral. In this paper we derive this collision integral from the equation of motion for the density correlation. By this way we are able to derive the spreading width within kinetic theory approach, from the coupling between one- and two-body dissipation. While the giant dipole and quadrupole resonances can be described within these models [19,20] the description of giant monopole resonances remains a problem [21–23]. The inclusion of memory effects does not improve the results much [22], but turns out to be necessary to obtain nonzero damping in the kinetic approach. In the latter paper it is shown that the consideration of quantal effects as well as finite size effects cannot account remarkably for this discrepancy. In [24,25] it was demonstrated that the contribution of collisional damping does not exceed 50% for the damping of giant dipole and 30% for giant monopole resonances. Therefore another damping mechanism should be responsible within the kinetic description of giant resonances. In this paper we show that this missing part is just the density-density fluctuations.

By a microscopic derivation of the collision integral including both density-density fluctuations as well as quasiparticle collisions we are able to rederive the former known

expression for the damping width, but with an additional enhancement factor. This factor connects the width of the resonances explicitly with the equation of state and in particular becomes large in the vicinity of a phase transition. In this way the kinetic transport equations are extended to include both the short and the long range correlations, which were separately considered in other recent approaches [26,27].

Next we consider the behavior of hot resonances. It is found that the recently observed saturation of damping with higher excitation energy is possible to explain within this treatment. The derived density dependence of the damping corresponding to the temperature behavior shows a transition from zero to first sound, which supports the result of [28]. We must remark that many studies have been devoted to the analysis of correlation effects in the theory of nuclear collective motions, in particular giant resonances [11,29,30]. It is however very difficult to predict the temperature dependence of these contributions. In this paper we suggest a semiclassical approach which also allows one to get a quite simple picture of correlation contributions at high temperature with clear consequences on the nature of collective motions in excited nuclear matter.

The outline of the paper is as follows. In Sec. II we give the derivation of the Lennard-Balescu collision integral which leads to an extension of common RPA to nonequilibrium situations, far from the near-equilibrium states, where linearization is justified. Together with the Boltzmann-Nordheim collision term describing short range correlations we derive a joint collision integral which combines both long range fluctuations and short range interactions. Here the complete non-Markovian behavior is considered. Then, in Sec. III, we give the calculation of the damping rate applying a linearization of the non-Markovian collision integrals due to external fields, but taking into account the memory effects. Analytic formulas are obtained, which include density-density fluctuations via an additional enhancement factor. The damping of isovector giant dipole and isoscalar monopole resonances is compared with the experimental values. A reasonable agreement is found. The saturation properties of giant resonances in hot nucleus are explained considering finite size effects and a transition from zero to first sound.

II. GENERAL COLLISION INTEGRAL

A. Density-density fluctuations

The density fluctuation in a nonequilibrium system is defined as

$$\delta\rho(11')=\Psi^+(1')\Psi(1)-\langle\Psi^+(1')\Psi(1)\rangle, \tag{1}$$

where numbers are cumulative variables $(r_1,t_1,s_1,i_1,\dots)$ denoting the space, time, spin, isospin, and other variables. It is convenient to describe fluctuations in the nonequilibrium many-particle system in terms of Green's functions. We introduce the correlated part of the two-particle Green's function including exchange by

$$L(121'2')=i[g_2(121'2')-g_1(11')g_1(22')]. \tag{2}$$

In the following we will consider the special correlation function

$$L(121^+2^+)=L(12) \tag{3}$$

which may be written as

$$L(12)=\Theta(t_1-t_2)L^>(12)+\Theta(t_2-t_1)L^<(12). \tag{4}$$

It follows immediately that the correlation function of density fluctuations is given by

$$\langle\delta\rho(11')\,\delta\rho(22')\rangle=iL^>(12)=iL^<(21). \tag{5}$$

We consider a system of fermions with the Hamiltonian

$$H=\sum_i\int d\mathbf{r}\Psi_i^*(\mathbf{r},t)H_i(\mathbf{r})\Psi_i(\mathbf{r},t)+\frac{1}{2}\sum_{i,j}\int d\overline{\mathbf{r}}d\mathbf{r}\Psi_i^*(\mathbf{r},t)\Psi_j^*(\overline{\mathbf{r}},t)V_{i,j}(\mathbf{r},\overline{\mathbf{r}})\Psi_i(\mathbf{r},t)\Psi_j(\overline{\mathbf{r}},t), \tag{6}$$

where V is a residual interaction.

With the help of the second equation of the Martin-Schwinger hierarchy, which couples the two-particle Green's function to the three-particle one, we get the following equation of motion for the correlator L defined in (2):

$$\begin{aligned}-i\left(i\hbar\frac{\partial}{\partial t_1}+\frac{(i\hbar\nabla_1)^2}{2m}\right)L(121'2')&=\delta(1-1')G(22')-\delta(1-2')G(21')-\delta(1-1')G(22')+G(22')\int d3V(13)L(131'3^+)\\&+\int d3V(13)G(33^+)L(121'2')+iG(11')G(22')\int d3V(13)G(33^+)\\&-i\int d3V(13)G_3(1231'2'3^+).\end{aligned} \tag{7}$$

Taking into account binary collisions allows one to describe short range interactions, while the long ranged interactions lead to characteristic collective modes. Approximations are necessary, which take into account collective effects. Such effects can be described in a convenient manner in terms of fluctuation quantities defined above. Whereas the binary collision approximation leads to the Boltzmann equation describing dilute systems, we want to include medium effects like collective excitations. Therefore we consider the three-particle Green's function in the following approximation [31]:

$$iG_3(1231'2'3^+)=iG(11')G(22')G(33^+)+G(11')L(232'3^+)+G(22')L(131'3^+)+G(33^+)L(121'2') \tag{8}$$

which describes the three-particle propagator in terms of all possible two-particle fluctuations. By this way we obtain an equation, which is equivalent to the RPA approximation in nonequilibrium. For Eq. (7) we now obtain [31]

$$\left(i\hbar\frac{\partial}{\partial t_1}+\frac{[(\hbar/i)\nabla_1]^2}{2m}\right)L(121'2')=-i\delta(1-2')G(21')-i\int d3V(13)L(232'3^+)G(11'). \tag{9}$$

The occurring integrations should now be specified and the appropriate solutions have to be selected. Therefore we rewrite the condition of the weakening of the initial correlation [31,32] by the definition of L, Eq. (2), into the following form:

$$L(121^+2^+)|_{t_1\to-\infty}=-iG(12^+)G(21^+)=L_0(121^+2^+). \tag{10}$$

Equation (9) for L can then be solved neglecting higher orders of interaction in the following manner:

$$L(421'2')=L_0(421'2')+L_0(411'1^+)V(13)L(232'3^+). \tag{11}$$

Because we are interested in the two-point-function, Eq. (3), we use a special case of Eq. (11):

$$L^{ab}(12)=L_0^{aa}(12)\delta^{ab}+\sum_c L_0^{aa}(14)V^{ac}(43)L^{cb}(32). \tag{12}$$

The different kinds of particles have been explicitly marked by Latin letters. Obviously the relation $L_0^{ab}\sim\delta^{ab}$ holds. In the following we will drop this notation for simplicity and restore these notations in the final results.

Equation (12) is a causal one, which means that all functions entering are causal ones. Therefore we can apply the Langreth-Wilkins [33,34] rules to obtain in operator notation

$$L^{\gtrless}=L_0^{\gtrless}+L_0^RVL^{\gtrless}+L_0^{\gtrless}VL^A,$$
$$L^{R/A}=L_0^{R/A}+L_0^{R/A}VL^{R/A}. \tag{13}$$

Using the second equation one gets from the first equation the optical theorem

$$L^{\gtrless}=(\epsilon^R)^{-1}L_0^{\gtrless}(\epsilon^A)^{-1}, \tag{14}$$

where we introduced the operators of retarded and advanced dielectric functions, which in the most general case are defined as (see, e.g. [35],)

$$\epsilon^R=1-L_0^RV\equiv(1+VL^R)^{-1},$$
$$\epsilon^A=1-VL_0^A\equiv(1+L^AV)^{-1}. \tag{15}$$

The optical theorem (14) is valid as long as initial fluctuations are sufficiently damped out and only the adiabatic evolution with time is considered [36]. This is certainly questionable for strong nonequilibrium and unstable situations. Therefore we will use another appropriate form later, which does not rely on the optical theorem.

With the help of these retarded and advanced response functions we can now write down the form of self-energy. From the definition of L in Eq. (2) and the introduction of self-energy [37], which is graphically represented in Fig. 1, we see that the following relation holds:

$$-\int d2V(12)L(121'2^+)=\int\Sigma'(12)G(21')d2. \tag{16}$$

Here Σ' denotes the self-energy with the Hartree part subtracted. The Fock part can further be split in order to get the correlated self-energy

$$\Sigma'(12)=\Sigma_c(12)\pm i\int d\overline{1}\overline{V}(1\overline{1})G^<(\overline{1}\overline{2})\delta(r_1-r_2),$$

where the upper sign stands for Bose particles and the lower one for Fermi particles. The required correlation self-energy is then easily derived in the following way:

$$\Sigma_{ab}^{\gtrless}(11')$$
$$=i\sum_c\int d2d2'V_{ab}(12)L_{bc}^{\gtrless}(22')V_{ca}(2'1')G_a^{\gtrless}(11'). \tag{17}$$

Furthermore, it is useful to introduce the fluctuating potential in the following way:

$$\sum_c V_{ac}(\epsilon_{cb}^a)^{-1}=\mathcal{V}_{ab}^a. \tag{18}$$

From (17) and (14) we have

$$\Sigma_{ab}^{\gtrless}(11')=i\sum_c\int d2d2'\mathcal{V}_{ab}^r(12)G_{bc}^{\gtrless}(2'2')$$
$$\times G_{bc}^{\lessgtr}(2'2')\mathcal{V}_{ca}^a(2'1')G_a^{\gtrless}(11'). \tag{19}$$

Another form of self-energy, without using the optical theorem, can be found from the diagrammatic expansion, Fig. 2, where the effective potential $\mathcal{V}$ sums all ring diagrams in Eq. (12):

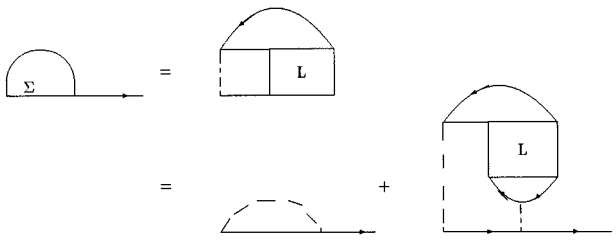

FIG. 1. Definition of self-energy in terms of density fluctuation.

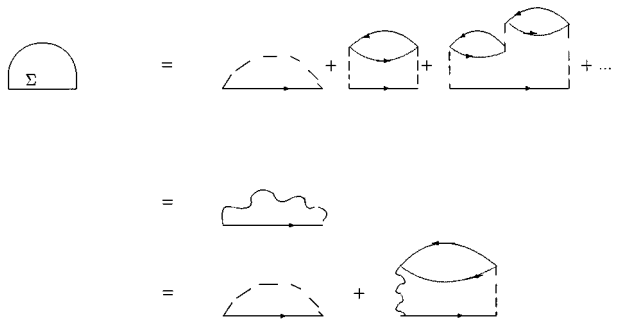

FIG. 2. Definition of the fluctuating potential.

$$\Sigma_{ab}(11')=i\sum_c \int d2d2' V_{ab}(12)G_{bc}(22') \times G_{bc}(22')\mathcal{V}_{ca}(2'1')G_a(11'). \quad (20)$$

In contrast to Eq. (19), now all functions appear to be causal ones. The corresponding correlation functions $\Sigma^{\gtrless}_{ab}$ can be derived easily with the help of the Lengreth-Wilkins rules. Applying the optical theorem (14) for $\mathcal{V}$ correspondingly one obtains just (19). Both forms will serve as a starting point to derive kinetic equations, but only (20) is appropriate for strong fluctuations, like those that can occur in the vicinity of phase transitions.

B. Lennard-Balescu collision integral with memory

In order to describe nuclear matter in nonequilibrium we consider now the general Kadanoff-Baym equation for the correlation function $g^{\gtrless}$ [37,38]:

$$i\left[\frac{\partial}{\partial t}+\left(\frac{p}{m_a}+\frac{\partial\Sigma^a_{HF}(pRt)}{\partial p}\right)\frac{\partial}{\partial R_a} - \frac{\partial\Sigma^a_{HF}(pRt)}{\partial R_a}\frac{\partial}{\partial p}\right]f^a_W(pRt)=I(p,t),$$

$$I(p,t)=\sum_b \int_0^\infty d\tau[\{g^>_a(p),\Sigma^<_{ab}(p)\} -\{g^<_a(p),\Sigma^>_{ab}(p)\}]_{(t-\tau/2,\tau)(t-\tau/2,-\tau)}. \quad (21)$$

Here $f_W(p,R,t)=\pm g^<(p,R,\tau=0,t)$ denotes the Wigner distribution function and $\{,\}$ is the anticommutator over integrals of Wigner coordinates indicated on the right. Here we restrict to the time diagonal case [34,39]. This equation is exact in time, but we used gradient expansion for space variables and drop all R dependence for simplicity.

Using the approximation for the self-energy, Eq. (17), and introducing the quasiparticle picture we derive the following collision integral [35]:

$$\overline{I^{\rm LB}_a}(k,t)=\mathrm{Re}\sum_b \int_0^t d\tau \int \frac{d\mathbf{k}'_a d\mathbf{k}_b d\mathbf{k}'_b}{(2\pi\hbar)^6}\delta(\mathbf{k}_a+\mathbf{k}_b-\mathbf{k}'_a-\mathbf{k}'_b)e^{-i[(E_{k_a}+E_{k_b}-E_{k'_a}-E_{k'_b})\tau]}[f_{a'}f_{b'}\overline{f_a f_b}-f_a f_b\overline{f_{a'}f_{b'}}]\mathcal{V}^R_{ab}(k'_b-k_b,E_{k'_b} -E_{k_b},t-\tau)\mathcal{V}^A_{ba}(k'_a-k_a,E_{k'_b}-E_{k_b},t) \quad (22)$$

with correspondingly $f_a=f(k_a,R,t-\tau)$, $\overline{f}=(1-f)$, and the quasiparticle energy E. If the latter one is time dependent, i.e., already in mean field approximation, then the exponential in (22) has to be replaced by

$$\exp\left(-i\int_t^{t-\tau}dt'[E_{k_a}(t')+E_{k_b}(t')-E_{k'_a}(t')-E_{k'_b}(t')]\right). \quad (23)$$

This collision integral is valid for quasiequilibrated fluctuations, which means that the initial fluctuations are already damped out [36]. For short time fluctuations, where the optical theorem is not applicable, Eq. (20) should be used as a starting point. At variance with (22) the fluctuating potential appears there only once. Here memory effects are incorporated. Equation (21) together with (22) determines the time evolution of the Wigner distribution function and therefore the behavior of the correlation function in the quasiparticle picture. The kinetic equation has the form of a generalized Lennard-Balescu (LB) equation and describes the influence of the density fluctuations on the collision process by a dynamical fluctuation potential.

C. Phase transition and critical regions

The kinetic equation (22) is a complicated coupled equation with the response function (15) given in the quasiparticle approximation as

$$\epsilon^R(p,\omega,R,t)=1+V(p)\int\frac{d\overline{\mathbf{p}}}{(2\pi\hbar)^3} \times\frac{f(E_{\overline{p-p}},R,t)-f(E_{\overline{p}},R,t)}{\omega-E_{\overline{p-p}}+E_{\overline{p}}+i\eta}. \quad (24)$$

With the help of this response function it is possible to find regions where the excitation spectrum is singular and therefore the density fluctuations become infinite. This is just the liquid gas phase transition region [40]. In Fig. 3 we plot the isothermal lines of the chemical potential. Here we use a Skyrme-type parametrization for the mean field [41]

$$U=A(n/n_0)+B(n/n_0)^\sigma \quad (25)$$

with

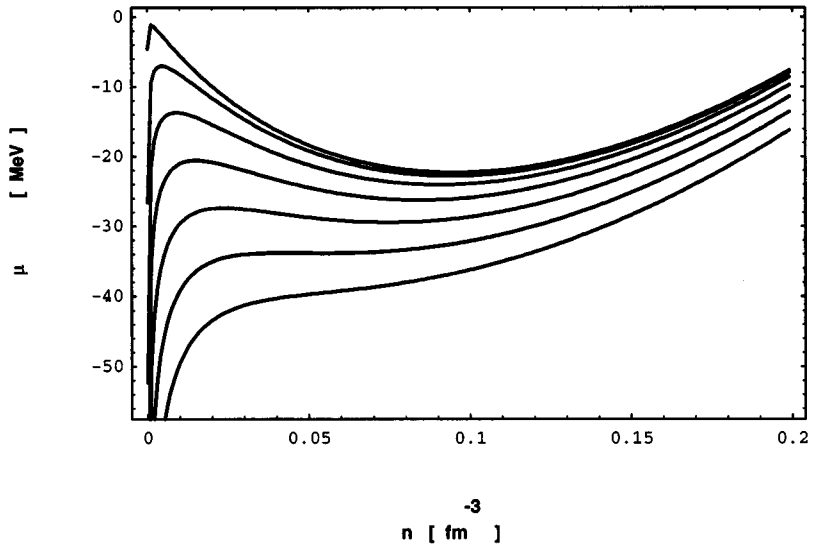

FIG. 3. Isothermal curves for chemical potential using the mean field parametrization (25). The temperatures are $T=1$, 4, 7, 10, 13, 16, 19 MeV.

$$\sigma=\frac{7}{6},\quad A=-356\ \text{MeV},\quad B=303\ \text{MeV},$$

$$n_0=0.145\ \text{fm}^{-3} \tag{26}$$

representing a soft equation of state.

D. Particle-particle scattering and BUU with memory

Up to now we have only considered the particle-hole channel. In order to include the particle-particle channel, which is necessary to describe short range interactions and repulsion, we use the T-matrix approximation. This leads to the collision integral [39]

$$I_a^{\mathrm{BUU}}(k_a,t)=\mathrm{Re}\sum_b\int\frac{d\mathbf{k}_a'd\mathbf{k}_b d\mathbf{k}_b'}{(2\pi\hbar)^6}\delta(\mathbf{k_a}+\mathbf{k_b}-\mathbf{k_a'}-\mathbf{k_b'})\int_0^t d\tau e(i/\hbar)(E_{k_a}+E_{k_b}-E_{k_a'}-E_{k_b'})\tau[f_{a'}f_{b'}\overline{f_a f_b}-f_a f_b\overline{f_{a'}}\,\overline{f_{b'}}]$$
$$\times\left\langle\frac{k_a-k_b}{2}\middle|T^R_{ab}(k_a+k_b,\epsilon_{k_a'}-\epsilon_{k_b'},t)\middle|\frac{k_a'-k_b'}{2}\right\rangle\left\langle\frac{k_a'-k_b'}{2}\middle|(T^A_{\mathrm{ex}})_{ab}(k_a+k_b,E_{k_a'}-E_{k_b'},t-\tau)\middle|\frac{k_a-k_b}{2}\right\rangle, \tag{27}$$

where $f_a=f(p,R,t-\tau)$. If the quasiparticle energy E becomes time dependent, the exponential in (27) has to be replaced by (23). Neglecting the retardation in the distribution functions and in the T matrix one obtains the usual quantum Boltzmann-Uehling-Uhlenbeck (BUU) collision integral equation.

E. Joint collision integral

Adding both collision integrals (22) and (27) we describe both the particle-particle collisions and the fluctuation or excitation properties. But if we look graphically at the sum of both expansions we see that the Fock part to the self-energy is double counted (see Fig. 4). This known fact [42] can be cured by subtracting just that Fock term. This translates into the subtraction of the Landau collision integral which is actually the Born term in the T-matrix equation. Thus, the final collision integral reads

$$I_{\mathrm{total}}=I_{\mathrm{BUU}}+I_{\mathrm{LB}}-I_{\mathrm{Landau}} \tag{28}$$

including both the scattering by BUU type of collision integrals and the fluctuating part by a Lennard-Balescu type of collision integral. It should be remarked that the ladder summation and the ring diagram summation can be derived on equal footing within the Faddeev equations [43].

In the following we proceed assuming small frequencies for the modes of the fluctuations. Further we use the quasiclassical limit or the long wavelength limit restricting to large scale fluctuations. Then the denominator of the fluctuating potential becomes [44]

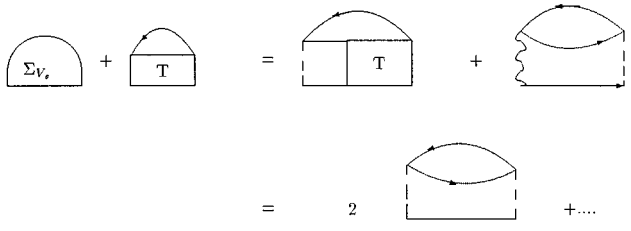

FIG. 4. Self-energy diagrams summed by the T matrix and the fluctuating potential. The Born term corresponding to the Landau collision integral is double counted and has to be subtracted once.

$$\epsilon(p,\omega)\to 1-V(0)\mathrm{Re}\Pi_0(0,0)=1+V(0)\frac{\partial n(\mu_0)}{\partial\mu_0} \quad (29)$$

with $V(0)=\partial U(n)/\partial n$ and the interaction-free chemical potential μ_0 . Using the effective potential (25) the denominator of the fluctuating potential can be written in front of the collision integral. Further, we assume that the scattering events may be well represented by a Born-type approximation. By this way the total collision integral, Eq. (28), reduces to the following form:

$$I_{\mathrm{joint}}(k_a,t)=\frac{1}{1+V(0)[\partial n(\mu_0)/\partial\mu_0]}\frac{1}{\hbar^2}\int\frac{d\overline{k_a}dk_b d\overline{k_b}}{(2\pi\hbar)^5}\delta^{(3)}(\overline{k_a}+\overline{k_b}-k_a-k_b)$$
$$\times\frac{|k_a-k_b|}{m}\frac{d\sigma}{d\Omega}\int_0^t d\tau\cos\left(\int_t^{t-\tau}(\overline{E_1}-E_1+E_{\overline{p}}-E_p)dt'\right)\{f_{a'}f_{b'}\overline{f_a}\overline{f_b}-f_af_b\overline{f_{a'}}\overline{f_{b'}}\} \quad (30)$$

with $f_a=f(k_a,R,t-\tau)$ and $\partial\sigma/\partial\Omega$ the differential cross section. We will refer to this collision integral as *joined collision integral*. It accounts for the contribution of the fluctuations as well as the particle-particle scattering. Further, the equation includes memory effects by the retardation of the distribution functions.

The factorization of the fluctuation denominator (29) out of the collision integral is only valid for soft modes $\omega=0$ and quasiclassical or long wavelengths $p=\hbar k\to 0$. Otherwise the scattering integral includes the full excitation function $\epsilon(p,\omega)$, which can be treated only numerically, which is in progress. This fluctuation factor is plotted in Fig. 5 for the mean field parametrization (25). One sees that in the vicinity of a phase transition region the factor becomes large. Inside the spinodal the factor is negative indicating the instability of the system. This is plotted in Fig. 6 where the fluctuation factor is shown for a temperature of $T=1$ MeV. While it remains almost a constant one for densities above 0.15 fm^{-3} it increases rapidly near the phase transition border around 0.11 fm^{-3}.

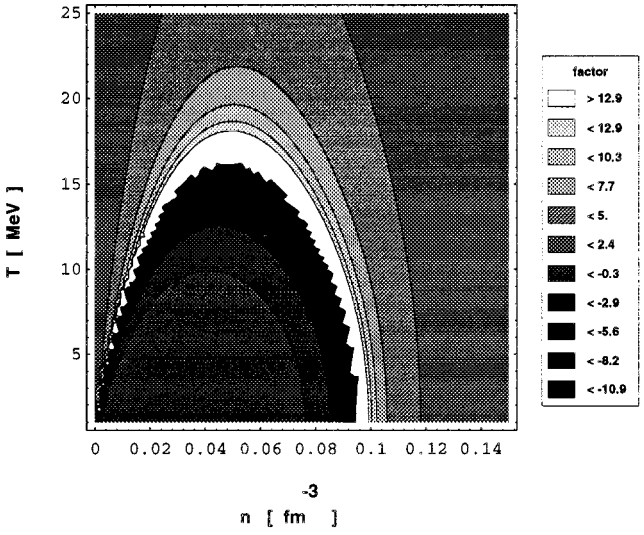

FIG. 5. The contour plot of the fluctuation enhancement factor vs density and temperature.

III. CALCULATION OF THE DAMPING RATE

The further considerations will be carried out in the same way as it was done in [23]. According to this model, the small deviations of the phase space density around equilibrium,

$$\delta f(r,p,t)=f(r,p,t)-f_0(\epsilon)=\chi(r,p,t)\frac{\partial}{\partial\epsilon}f_0, \quad (31)$$

obey the linearized transport equation

$$\frac{\partial}{\partial t}\delta f-\{\delta h,f_0\}-\{h_0,\delta f\}$$
$$=\int dp_2dp_3dp_4\delta W[(1-f_0)(1-f_2^0)f_3^0f_4^0$$
$$-f_0f_2^0(1-f_3^0)(1-f_4^0)], \quad (32)$$

where f_0 is the Fermi function and the transition rate is given by

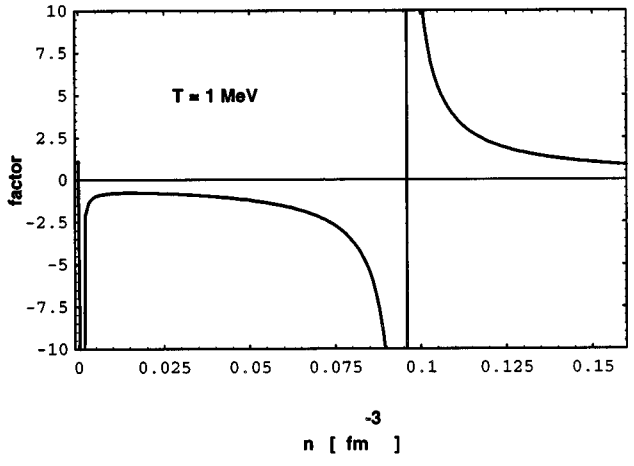

FIG. 6. The fluctuation enhancement factor of Fig. 5 for a temperature of 1 MeV vs density.

$$\delta W = \frac{W}{1+V(0)[\partial n(\mu_0)/\partial\mu_0]} \int \frac{d\omega}{2\pi} e^{-i\omega t} \Delta\chi(\omega) \frac{Z(\omega)}{2}, \tag{33}$$

with $Z(\omega)=[\delta(\Delta\epsilon-\omega)-\delta(\Delta\epsilon+\omega)]/\omega$, $\Delta\epsilon=\epsilon_3+\epsilon_4-\epsilon_1-\epsilon_2$, and $\Delta\chi=\chi_3+\chi_4-\chi_1-\chi_2$. Here we extend the result [23] since the BUU transition rate W term is divided by a factor due to the fluctuating term (29). We see that in the vicinity of phase transition we have an overall enhancement of collisions. This critical scattering will be published elsewhere [40].

Now we turn to the evaluation of the extended collisional damping for giant modes. The relaxation rate for isoscalar harmonic vibrations with a mean frequency Ω can be defined as [23]

$$\frac{1}{\tau} = -\frac{\int dr dp \chi \delta K}{\int dr dp \chi \delta f} = \frac{\int dr dp_1 dp_2 dp_3 dp_4 WZ(\Omega)(\Delta\chi)^2 f_1^0 f_2^0 (1-f_3^0)(1-f_4^0)}{\int dr dp \chi^2 (\partial/\partial\epsilon) f_0}. \tag{34}$$

Similar expressions have been derived for the rates of isovector vibrations [23]. There the distortion of the momentum distribution is used in terms of the Legendre functions as $\chi=p^2P_2(\Theta)$ for quadrupole vibrations and $\chi=pP_1(\Theta)$ for giant dipole vibrations.

A. Spreading width of giant monopole resonances

Within the context of the scaling model description of monopole mode, the Fermi surface distortion is assumed simplified to be $\chi=p^2$ and the damping of giant monopole modes $\Gamma=\hbar/\tau$ is given according to [22,23] by

$$\Gamma = \frac{\hbar v_F \sigma n}{1+V(n)(\partial n/\partial\mu_0)} \left(\frac{\hbar\Omega}{4\epsilon_F}\right)^2 \left[\left(\frac{\hbar\Omega}{\epsilon_F}\right)^2 + \left(\frac{2\pi T}{\epsilon_F}\right)^2\right], \tag{35}$$

where temperature is assumed to be small compared with the Fermi energy so that standard integration techniques have been applied [45]. If we had neglected the memory effects we would have obtained no damping of the monopole resonances at all. This underlines the importance of incorporating memory effects for the description of giant monopole resonances. A more detailed analysis of the distortion of the Fermi surface can be found in [21], where the effect of the nonspherical Fermi surface is considered. This would lead to a minor correction of the presented results. We restrict here to the main effect of fluctuations and assume therefore a simplified scaling model for the distortion. The numerical treatment of (34) in comparison with (35) as well as the inclusion of diffusive surface can be found in [24], which leads to an enhancement but still underestimates the experimental values.

The expression, Eq. (35), is the generalization of the damping rates derived previously [22] by incorporating the density fluctuations. These fluctuations yield a direct coupling of the damping rate to the equation of state.

For comparison with the experiment we used the nuclear matter parameters as $n=0.16$ fm^{-3}, $v_F=0.28c$, $\epsilon_F=37$ MeV, the cross section $\sigma=40$ mb, and the mass dependence of the resonance energies was assumed to be the conventional one [7,8].

The results for the giant monopole damping versus mass number in the ground state ($T=0$) can be found in Fig. 7 where we plot Eq. (35) with and without the fluctuation enhancement factor. While the lower curve represents the damping rate without fluctuations and underestimates the experimental values considerably, the upper curve is the spreading width including fluctuations. The density fluctuations lead to an enhancement due to the forefactor. This enhancement factor is explicitly seen by applying the Sommerfeld expansion

$$1+V(n)\frac{\partial n}{\partial\mu_0} = 1+V(n)\frac{mp_f}{\pi^2\hbar^3}\left[1-\frac{1}{48}\left(\frac{2\pi T}{\epsilon_F}\right)^2\right], \tag{36}$$

where μ_0 denotes the free chemical potential. One recognizes that besides the quadratic increase of the damping width with temperature there is an additional increase by this enhancement factor. Here we have assumed an effective density for the fluctuations of 0.107 fm^{-3}.

This means that it is possible to reproduce the monopole damping width by considering density fluctuations and assuming a lower density than nuclear matter density. This can be related to finite nuclear size effects as follows. We ask in a first step which density has to be assumed to reproduce the damping including fluctuation effects. With the help of the Woods-Saxon density profile for the finite nucleus

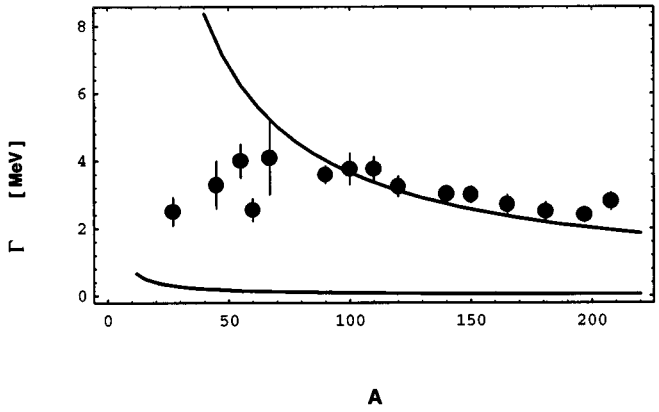

FIG. 7. The experimental values of the damping width of giant monopole resonances together with the theoretical damping according to (35) with (upper) and without (lower curve) fluctuation.

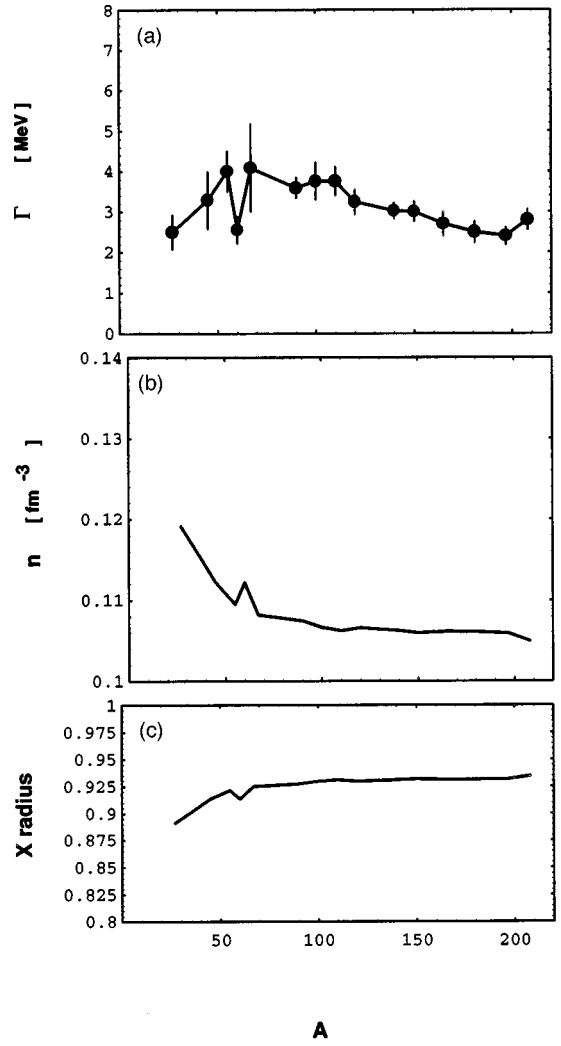

FIG. 8. The fit of the density (b), where the giant monopole resonances are built up (a), together with the corresponding fractional radius $r=X\times R_m$ where this density is realized (c). The density profile is assumed via a Woods-Saxon potential with radius R_m.

$$n(r)=\frac{n_0}{e^{(r-R_m)/a}+1}, \tag{37}$$

we can fit the effective radius $r=x\times R_m$ such that the corresponding density reproduces the experimental values. This can be seen in Fig. 8. The fitted density to reproduce the experimental values is plotted in Fig. 8(b). In Fig. 8(c) the factor x to be multiplied with the Woods-Saxon radius R_m of the finite nucleus is shown. One recognizes that we can reproduce the width of the giant resonances assuming that the resonance affects the nucleus up to a range of $r=0.92\times R_m$. This means that the surface region around $r=0.92\times R_m$ is the most important for density fluctuations and therefore for the damping rate of the resonances. On the other hand it is reasonable to have enhanced density fluctuations on the surface just because we expect larger relative variances where the average occupation is smaller. However we have to remark that in this approach the Landau damping is still missing and thus probably we have to use a quite large enhancement factor, i.e., a too small effective density for the evaluation of fluctuations.

B. Giant dipole resonances

In analogy to the derivation of the formula, Eq. (35), the width of the giant dipole resonances (GDR's) can be given by

$$\Gamma_{\mathrm{GDR}}=\frac{\hbar}{8}\frac{v_F\sigma_v n}{1+V(n)(\partial n/\partial\mu_0)}\left[\left(\frac{\hbar\Omega}{\epsilon_F}\right)^2+\left(\frac{2\pi T}{\epsilon_F}\right)^2\right], \tag{38}$$

where $\sigma_v=(1/2)\sigma_{pn}$ and we have in generalization to [23] the fluctuation enhancement factor. This damping rate is plotted with and without fluctuations together with the experimental values in Fig. 9. Here we used an effective density of 0.122 fm^{-3} for the fluctuation to reproduce the data in the upper curve. As before, in the case of monopole resonances, we can reproduce the damping width assuming a smaller density for the fluctuations. Apart from the Landau damping, as mentioned before, in this case we have also deformation effects present in the data.

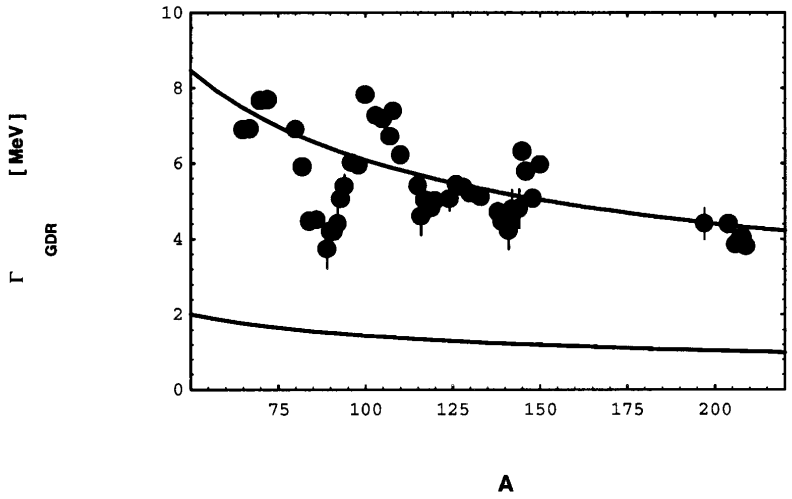

FIG. 9. The experimental values of the damping width of giant dipole resonances together with the theoretical damping according to (35) with (upper) and without fluctuation (lower curve).

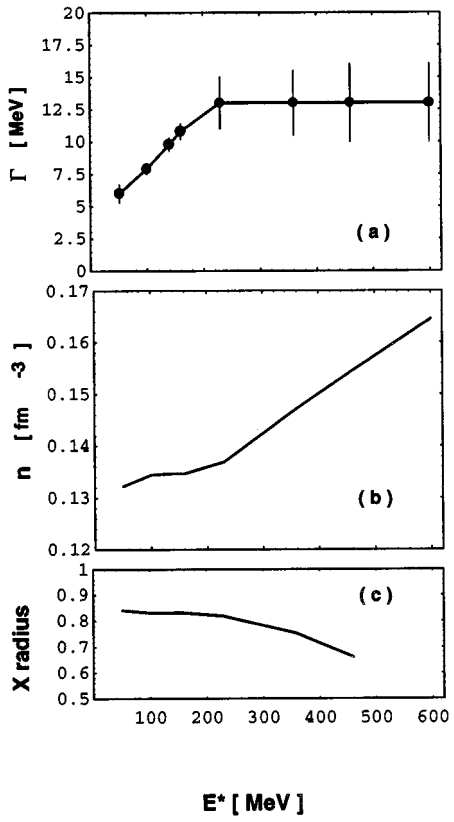

FIG. 10. (a) Experimental widths of giant dipole resonances for $^{108-112}$ Sn vs excitation energy. The data are taken from [46,56–59,47,6]. (b) Fit of the effective density, as explained in the text. (c) Corresponding fraction of radius $r=X\times R_m$ where this density is realized. The saturation of damping width is connected with an increasing size of the nucleus which is affected by density oscillations.

A comparison with the monopole damping shows that the effective surface is about $0.86R_m$ for dipole vibrations while it was $0.92R_m$ for monopole vibrations. This seems to indicate that the higher the order of multipolarity the less important the fluctuation effect.

C. Saturation of damping for hot nucleus

It is observed [46–50] that the total width of giant dipole resonances seems to saturate in heated nuclei. Furthermore, a complete vanishing at higher excitation energies above 600 MeV is observed. The most investigated nuclei at high excitation energies are $^{108-112}$Sn. While the centroid energy remains nearly constant the width increases with excitation energy and can be fitted by $\Gamma=4.8+0.0026E_{\rm cn}^{1.6}$ in the range of low energies [51,48,49]. At higher energies the total width appears to show a saturation above 200 MeV.

In our model we can easily get a saturation for the spreading width, with some interesting implications. We assume a functional dependence of the excitation energy from the temperature via $E_{\rm cn}=(\pi^2A/4\epsilon_F)T^2$. The temperature dependence of the centroid energy of the giant resonances can be neglected [52]. Then in Fig. 10 we fit the effective fluctuation density corresponding to our general formula (38) to the experimental values. From the result in Fig. 10(b) one sees that the saturation is connected to an increase of the effective density. This corresponds to a larger effective surface region for the resonance, as represented in Fig. 10(c). We always have to remind the reader that Landau damping effects are not included in this calculation and so we are probably overestimating the effects of fluctuations.

We also see that with increasing excitation energy the effective surface range which is important for the giant dipole resonances becomes larger and reaches almost half of the nucleus radius at about 500 MeV. In other words, for higher excitation energies the finite nucleus will be simply destroyed by too strong density fluctuations.

Now it is important to understand the physical mechanism which leads to a transient regime around 200 MeV, where the saturation occurs. While the sharp increase for smaller temperatures is explained by a typical zero sound attenuation due to the effect of collisions, although incorporating density-density fluctuations, the saturation requires a new understanding. In [28] it was pointed out that this saturation can be due to the transition between zero and first sound. This observation is underlined by our calculation which clearly shows how density fluctuations are providing a substantial contribution to the transition to first sound propagation. Indeed if we use the expression, Eq. (38), *without the fluctuation correction* for the collision time τ we can easily show that the Landau condition $\Omega\tau\simeq 1$ [53,54] for that transition can be only reached at quite high temperature, of the order of 12–15 MeV, much more than an equilibrated compound nucleus can sustain [55]. This is also valid because the temperature dependence of the resonance centroid energy can be neglected [55,52]. The presence of the reduction factor, Eq. (36), in the collision time drastically changes the scenario in the direction of a strong reduction of the transition temperature. Using the same effective density for the evaluation of fluctuation effects as in the ground state case, we reach the transition condition at temperatures of the order of 4 MeV, i.e., at excitation energies around 200 MeV. This means that at higher excitation energy the propagation of the dipole vibration is of classical type, with a related decreasing temperature behavior of the attenuation. We must remark that for heated dipole states we have always some evaporation width which is increasing with temperature. However the zero first sound transition will open the possibility of observing collective dipole emissions up to very hot nuclei, near the limiting temperature for a formation of a compound nucleus.

IV. SUMMARY

Within the two time Green's functions formalism the equation of motion for the density-density fluctuations is derived. The RPA leads to a generalized kinetic equation of Lennard-Balescu type. A joint kinetic equation is presented which accounts for the long range density-density fluctuations and for the short range correlations by particle-particle scattering. This kinetic equation includes memory effects via retardation of the distribution functions and off shell effects by explicit time dependence of the transition matrix element.

Within the soft mode and small wave vector regime an approximative kinetic equation is derived. Here the fluctua-

tions lead to a forefactor of the collision integral. The memory effect is included within retardation of the single-particle distribution function. Assuming scaling forms for the distortion of the Fermi sphere the expressions for the damping rates of giant monopole and dipole resonances are derived. The density fluctuations lead to an additional forefactor coupling the damping rate explicitly to the equation of state. By this way we can show an overall enhancement factor for the theoretical kinetic damping rate which reproduces the experimental values. It is shown that the fluctuation effects become more important for the damping for lower multipolarity. The damping of monopole resonances can be well reproduced up to nuclei with mass number 50, although the Landau single-particle damping is not included here.

The problem of saturation of the giant dipole width at high excitation energies is explained by a transition from zero to first sound in the collective mode propagation. The inclusion of fluctuations is essential in order to have this transition at relatively low temperatures in excited nuclear matter. This seems to be in agreement with recent observations and could open new exciting perspectives in theory as well as in experiments.

ACKNOWLEDGMENTS

The authors thank S. Ayik for interesting discussions and Armen Sedrakian for reading the manuscript and useful comments. The authors also thank the late Ludwig Münchow, whose inspiring ideas have greatly contributed to this work. This work was supported by a grant (K.M.) of the German Academic Exchange Service and the hospitality of Tennessee Technological University, where part of this work was carried out, is gratefully acknowledged.

[1] G. G. Baldwin and G. S. Klaiber, Phys. Rev. **71**, 3 (1947).
[2] B. L. Berman, Nucl. Data Tables **15**, 321 (1975).
[3] S. S. Dietrich and B. L. Berman, Atomic Data Nucl. Data Tables **38**, 199 (1988).
[4] H. Emling, Prog. Part. Nucl. Phys. **33**, 729 (1994).
[5] B. L. Berman and S. C. Fultz, Rev. Mod. Phys. **47**, 713 (1975).
[6] J. J. Gaardhoje, in *Frontiers in Nuclear Dynamics*, edited by R. A. Broglia and C. H. Dasso (Plenum, New York, 1992), Vol. 425, p. 133.
[7] A. van der Woude, Prog. Part. Nucl. Phys. **18**, 217 (1987).
[8] A. van der Woude, in *International Review of Nuclear Physics*, edited by J. Speth (World Scientific, Singapore, 1991), Vol. 7, p. 99.
[9] P.-G. Reinhard, S. Weisgerber, O. Genzken, and M. Brack, Z. Phys. A **349**, 219 (1994).
[10] G. F. Bertsch, P. F. Bortignon, and R. A. Broglia, Rev. Mod. Phys. **55**, 287 (1983).
[11] J. Speth and J. Wambach, in *International Review of Nuclear Physics* [8], p. 1.
[12] G. F. Bertsch and S. D. Gupta, Phys. Rep. **160**, 189 (1988).
[13] C. Yanhuang and M. DiToro, Phys. Rev. C **39**, 105 (1989).
[14] A. Bonasera, G. F. Burgio, M. DiToro, and H. H. Wolter, Phys. Rev. C **39**, 2385 (1989).
[15] M. DiToro and C. Yanhuang, Phys. Rev. C **42**, 2748 (1990).
[16] D. Brink, Nucl. Phys. A **482**, 3 (1988).
[17] D. Brink, A. Dellafiore, and M. DiToro, Nucl. Phys. A **456**, 205 (1986).
[18] A. Dellafiore and F. Matera, Nucl. Phys. A **460**, 265 (1986).
[19] I. N. Mikhalov, J. Piperova, M. DiToro, and A. Smerzi, Nucl. Phys. A **575**, 269 (1994).
[20] A. Bonasera, M. DiToro, and F. Gulminelli, Phys. Rev. C **42**, 966 (1990).
[21] V. Abrosimov, M. DiToro, and A. Smerzi, Z. Phys. A **347**, 161 (1994).
[22] S. Ayik, M. Belkacem, and A. Bonasera, Phys. Rev. C **51**, 611 (1995).
[23] S. Ayik and D. Boilley, Phys. Lett. B **276**, 263 (1992); **284**, 482 (1992).
[24] M. Belkacem, S. Ayik, and A. Bonasera, Phys. Rev. C **52**, 2499 (1995).
[25] V. M. Kolomietz, V. A. Pluiko, and S. Shlomo, Phys. Rev. C **52**, 2480 (1995).
[26] M. Colonna, Ph. Chomaz, and J. Randrup, Nucl. Phys. A **567**, 637 (1994).
[27] V. N. Kondratyev and M. DiToro, Phys. Rev. C **53**, 5 (1996).
[28] V. Baran *et al.*, Nucl. Phys. **A599**, 29c (1996).
[29] G. F. Bertsch, P. F. Bortignon, and R. A. Broglia, Rev. Mod. Phys. **55**, 287 (1989).
[30] V. G. Soloviev, Z. Phys. A **334**, 143 (1989); Prog. Part. Nucl. Phys. **28**, 49 (1992), and references therein.
[31] M. Schlanges and T. Bornath, Wiss. Z. Universität Rostock **36**, 65 (1987).
[32] Y. L. Klimontovich, D. Kremp, and W. D. Kraeft, *Advances Chemistry and Physics* (Wiley, New York, 1987).
[33] D. C. Langreth and J. W. Wilkins, Phys. Rev. B **6**, 3189 (1972).
[34] A. P. Jauho and J. W. Wilkins, Phys. Rev. B **29**, 1919 (1984).
[35] K. Morawetz, Phys. Rev. E **50**, 4625 (1994).
[36] B. Bezzerides and D. F. DuBois, Ann. Phys. (N.Y.) **70**, 10 (1972).
[37] L. P. Kadanoff and G. Baym, *Quantum Statistical Mechanics* (Benjamin, New York, 1962).
[38] L. V. Keldysh, Zh. Éksp. Teor. Fiz. **47**, 1515 (1964) [Sov. Phys. JETP **20**, 1018 (1965)].
[39] K. Morawetz and G. Röpke, Phys. Rev. E **51**, 4246 (1995).
[40] L. Münchow, K. Morawetz, G. Röpke, and M. DiToro, submitted to Nucl. Phys. A.
[41] Y. Ivanov and S. Ayik (unpublished).
[42] P. Danielewicz, Ann. Phys. (N.Y.) **152**, 239 (1984).
[43] P. Danielewicz, Ann. Phys. (N.Y.) **197**, 154 (1990).
[44] K. Morawetz and G. Roepke, in *Proceedings of the International Workshop* (GSI, Hirschegg, Kleinwalsertal, 1993), p. 292.
[45] H. Smith and H. Hojgaard-Jensen, *Transport Phenomena* (Clarendon Press, Oxford, 1989).
[46] J. J. Gaardhoje, in *Annual Review of Nuclear Particles in Science*, edited by J. D. Jackson (Palo Alto, California, 1992), Vol. 42, p. 483.
[47] A. Bracco *et. al.*, Phys. Rev. Lett. **62**, 2080 (1989).

[48] K. Yoshida *et. al.*, Phys. Lett. B **245**, 7 (1990).
[49] J. Kasagi *et al.*, Nucl. Phys. A **538**, 585 (1992).
[50] T. Suomijaervi *et al.*, Nucl. Phys. A **583**, 105c (1994).
[51] D. R. Chakrabarty *et al.*, Phys. Rev. C **36**, 1886 (1987).
[52] W. Wen, G. Dai, and G. Jin, Phys. Rev. C **52**, 187 (1995).
[53] L. D. Landau, Sov. Phys. JETP **5**, 101 (1957).
[54] E. Lifschitz and L. P. Pitaevsky, in *Theoretical Physics*, edited by E. Lifschitz (Akademie Verlag, Berlin, 1981).
[55] M. Colonna, M. DiToro, V. Latora, and A. Smerzi, Prog. Part. Nucl. Phys. **30**, 17 (1993).
[56] J. Gaardhoje, C. Ellegaard, B. Herskind, and S. G. Steadman, Phys. Rev. Lett. **53**, 148 (1984).
[57] J. Gaardhoje *et al.*, Phys. Rev. Lett. **56**, 1783 (1986).
[58] J. Gaardhoje *et al.*, Phys. Rev. Lett. **59**, 1409 (1987).
[59] D. R. Chakrabarty *et al.*, Phys. Rev. C **36**, 1886 (1987).