Interplay between one-body and collisional damping of collective motion in nuclei

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Damping of giant collective vibrations in nuclei is studied within the framework of the Landau-Vlasov kinetic equation. A phenomenological method of independent sources of dissipation is proposed for taking into account the contributions of one-body dissipation, the relaxation due to the two-body collisions and the particle emission. An expression for the intrinsic width of slow damped collective vibrations is obtained. In the general case, this expression cannot be represented as a sum of the widths associated with the different independent sources of the damping. This is a peculiarity of the collisional Landau-Vlasov equation where the Fermi-surface distortion effect influences both the self-consistent mean field and the memory effect at the relaxation processes. The interplay between the one-body, the two-body, and the particle emission channels which contribute to the formation of the total intrinsic width of the isoscalar 2^+ and 3^- and isovector 1^- giant multipole resonances in cold and hot nuclei is discussed. We have shown that the criterion for the transition temperature T_{tr} between the zero-sound and first-sound regimes in hot nuclei is different from the case of infinite nuclear matter due to the contribution from the one-body relaxation and the particle emission. In the case of the isovector GDR the corresponding transition can be reached at temperature $T_{\mathrm{tr}}=4-5$ MeV. [S0556-2813(96)05112-6]

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I. INTRODUCTION

The interplay between different relaxation mechanisms of collective motion and their dependence on the temperature in many-body systems is a subject of wide investigation [1–6]. In the present paper we consider the damping of the nuclear multipole vibrations within the semiclassical Landau-Vlasov kinetic theory. Semiclassical methods seem to be quite instructive for an investigation of the averaged properties of the multiparticle systems. In many cases, they allow us to obtain analytical results and represent them in a transparent way.

In what follows we concentrate on the investigation of the contributions of different relaxation mechanisms to the intrinsic width of the giant multipole resonance (GMR). We determine the intrinsic width as formed by three main sources: (i) The relaxation due to the coupling of both particle and hole to more complicated states lying at the same excitation energy. This is the so-called two-body collisional damping. In the kinetic theory, this type of relaxation is simulated by the collision integral and leads to the collisional component of the intrinsic width. (ii) The fragmentation width caused by the interaction of particles with the time-dependent self-consistent mean field. In the quantum random phase approximation (RPA) calculations this contribution to the width does not reflect a motion of system towards the thermal equilibrium but indicates rather a redistribution of the particle-hole excitations in a vicinity of the collective state. In our kinetic approach, we will imitate the fragmentation width by the one-body (''wall'') relaxation. In agreement with the above-mentioned peculiarity of the fragmentation width, the one-body relaxation is conceived as a dissipative phenomenon only if the observation time is shorter than the Poincare time. Note also that it was shown in Refs. [7–9], in the classical limit for the random phase approximation, the fragmentation width coincides with the width obtained from the one-body relaxation mechanism. (iii) The emission of the particles into the free space.

To account for these three relaxation mechanisms in the evaluation of the GMR width we propose in Sec. II a phenomenological approach. We take into account the two-body (collisional) damping exactly, incorporating into the collision integral the memory effects associated with the mean-field vibrations. The one-body relaxation and the coupling to the continuum states (particle's emission) are taken into consideration approximately by adding to the Landau-Vlasov equation some source terms. These source terms are assumed to be independent of each other. They characterize relaxation in a given channel only in the absence of all other damping mechanisms. All source terms are included in the form of the relaxation time approximation. The total change rate in the distribution function is taken as a sum of the change rates in various damping channels (independent dissipation rates approximation).

In Sec. III the relaxation times in one-body and two-body damping channels are analyzed and the intrinsic width of the GMR in cold and hot nuclei is calculated. We also discuss the interplay between different damping channels which contribute to the formation of the total intrinsic widths of the GMR. The numerical results and general discussion of the mass number and temperature dependences of the different relaxation mechanisms are presented in Sec. IV. The conclusion and summary are given in Sec. V.

II. WIDTHS OF GMR WITHIN THE INDEPENDENT DISSIPATION RATES APPROXIMATION

We will start from the Landau-Vlasov equation for the phase space distribution function $f(\mathbf{r},\mathbf{p},t)$, completed by a source term $J(\{f\})$ for relaxation processes. The two-component nuclear Fermi system has both the isoscalar and isovector excitations. In quantum calculations these excitations are distinguished by the isospin quantum number. In macroscopic approaches (see, for example, Refs. [10–13]) the isoscalar and isovector modes correspond to the in-phase and out-of-phase motions of neutrons and protons, respectively. That means that both modes can be described in terms of the distortions of the distribution function in the form $\delta f^{(\pm)} = \delta f_p \pm \delta f_n$ with $\delta f_p = \pm \delta f_n$, where the subindices p or n label protons or neutrons and the plus or minus sign denotes the isoscalar or isovector modes, respectively. Neglecting a small difference in the chemical potentials for protons and neutrons and assuming $f_{0,p} = f_{0,n} = f_0$, where $f_0 \equiv f_0(\mathbf{r},\mathbf{p})$ is the equilibrium distribution function, we write down the linearized two-component Landau-Vlasov equation in the form

$$\frac{\partial \delta f^{(\pm)}}{\partial t} + \frac{\mathbf{p}}{m}\frac{\partial \delta f^{(\pm)}}{\partial \mathbf{r}} - \frac{\partial V_0}{\partial \mathbf{r}}\frac{\partial \delta f^{(\pm)}}{\partial \mathbf{p}} - \frac{\partial \delta V^{(\pm)}}{\partial \mathbf{r}}\frac{\partial f_0}{\partial \mathbf{p}} = J^{(\pm)}(\{\delta f_p, \delta f_n\}), \quad (1)$$

where $\delta V^{(\pm)} \equiv \delta V^{(\pm)}(\mathbf{r},\mathbf{p},t)$ is the Wigner transform of the variation of the self-consistent potential with respect to the equilibrium value V_0. In the nuclear interior the mean field variation $\delta V^{(\pm)}$ can be expressed in terms of the Landau interaction amplitude $F^{(\pm)}(\mathbf{p},\mathbf{p}')$ as

$$\delta V^{(\pm)} = \frac{g}{N_F}\int \frac{d\mathbf{p}'}{(2\pi\hbar)^3} F^{(\pm)}(\mathbf{p},\mathbf{p}')\,\delta f^{(\pm)}(\mathbf{r},\mathbf{p}';t), \quad (2)$$

where $N_F = 2p_F m^*/(g\pi^2\hbar^3)$, p_F is the Fermi momentum, m^* is the effective mass of nucleon, and g is the spin degeneracy factor. The quantity $F^{(\pm)}(\mathbf{p},\mathbf{p}')$ is usually parameterized in terms of the Landau constants $F_0^{(\pm)}$ and $F_1^{(\pm)}$ as

$$F^{(\pm)}(\mathbf{p},\mathbf{p}') = F_0^{(\pm)} + F_1^{(\pm)}(\hat{p}\cdot\hat{p}'). \quad (3)$$

In the isoscalar case, the Landau constants are related to the incompressibility modulus $K \simeq 220$ MeV [14] of matter and the effective mass $m^* \simeq 0.8m$ [15] by

$$K = 6\mu(1 + F_0^{(+)}), \quad m^* = m(1 + F_1^{(+)}/3). \quad (4)$$

Here m is the mass of free nucleon and μ is the chemical potential. We have that $\mu \approx \epsilon_F = p_F^2/2m^*$ for $T \ll \epsilon_F$, where ϵ_F is the Fermi energy and T is the temperature. In the isovector case, the Landau parameter $F_0^{(-)}$ is related to the nuclear symmetry energy $b_{\text{symm}} \simeq 30$ MeV [10]. Namely [16],

$$b_{\text{symm}} = \frac{1}{3}\mu(1 + F_0^{(-)}). \quad (5)$$

To simplify the presentation, we will omit in the following the superscripts $(\pm)$ and include them only when it is necessary to avoid confusion.

The right-hand side of Eq. (1) represents the change of the distribution function due to relaxation. In this work we use the approximation of independent dissipation rates. Namely, we assume

$$J(\{\delta f\}) = J_c(\{\delta f\}) + J_s(\{\delta f\}) + J_{\uparrow}(\{\delta f\}), \quad (6)$$

where $J_c(\{\delta f\})$ is the collision integral for the two-body collisions, $J_s(\{\delta f\})$ determines the change in the distribution function resulting from one-body relaxation, and $J_{\uparrow}(\{\delta f\})$ takes into account the possibility of particle emission. Strictly speaking, the relation (6) is valid only approximately since it suggests that any relaxation mechanism, in the absence of the other ones, drives toward the same final distribution function.

We now comment on the one-body relaxation and the possibility of treatment of this one as a source term in the kinetic equation. In a system like a nucleus, where the self-consistent mean field decreases sharply in a small region of the coordinate space in comparison with the bulk dimensions of the system, one can define a surface. In this case, the problem of solving the kinetic equation in coordinate space is equivalent to a boundary value problem with some boundary conditions on the surface [17,18]. Considering this surface (more precisely the parameters describing it) as a collective variable, the process of energy exchange between the inner (particles) and collective (surface) degrees of freedom can be described as a relaxation of the collective motion. In this sense we can talk about the presence of relaxation even in the collisionless Landau-Vlasov equation. The origin of this relaxation is the collision of the particles with a moving surface and this type of relaxation is related to the fragmentation width of the collective states in quantum calculations like the RPA (see corresponding discussion in previous section). In fact, this type of relaxation in a transport theory is governed by the boundary conditions imposed on the surface of the system [19]. For instance, the wall formula [20] for dissipation energy in a semi-infinite Fermi liquid results from the boundary conditions of the specular reflection of particles from the moving surface [17,18,21–23]. Below we will consider only this variety of the one-body relaxation and call it one-body wall dissipation. It was shown in Refs. [24,25] that the problem of finding a solution of the kinetic equation with the boundary conditions of the specular or diffuse reflections is identical to the problem of finding a solution of the kinetic equation with source terms. Thus, at least for the boundary conditions of the specular and diffuse reflections one can simulate the one-body wall dissipation by a source term as it was assumed in Eqs. (1) and (6).

We will consider the terms $J_s(\{\delta f\})$ and $J_{\uparrow}(\{\delta f\})$ in Eq. (6) within the relaxation time approximation of the form

$$J_s(\{\delta f\}) = -\frac{\delta f(\mathbf{r},\mathbf{p},t)}{\tau_s}, \quad J_{\uparrow}(\{\delta f\}) = -\frac{\delta f(\mathbf{r},\mathbf{p},t)}{\tau_e}. \quad (7)$$

Here, τ_s is the relaxation time corresponding to the equilibration of the system due to the one-body wall dissipation. The particle emission relaxation time τ_e is determined by the life time of the collective excitation with respect to an emission of particles into the continuum. The choice of both relaxation times will be discussed later.

We point out that the one-body source term $J_s(\{\delta f\})$ does not contain the components with multipolarities $\ell = 0$ and $\ell = 1$ for the distorted distribution function δf in momentum space because of the conservation of the number of particles and the total momentum. In the case of the particle emission term $J_{\uparrow}$, we neglect the components with $\ell=0$ and 1 expecting that, at not too high temperatures, the emission of particles does not change essentially the mean particle density and the current density in the nuclear interior. Note also that the collision integral in the relaxation time approximation for one-body dissipation was used earlier in Refs. [26,27].

For small amplitude eigenvibrations and small deviations of the momentum distribution from the Fermi sphere, the collision integral $J_c(\{\delta f\})$ in Eq. (6) can be linearized in δf and represented in the form of the generalized τ approximation [6,11,28–31]

$$J_c(\{\delta f\}) = -\frac{\delta f(\mathbf{r},\mathbf{p},t)}{\tau_c(\omega_0,T)}, \tag{8}$$

where $\tau_c(\omega_0,T)$ is the relaxation time due to the interparticle collisions and ω_0 is the eigenfrequency of the collective eigenvibrations. The ω dependence in Eq. (8) takes into account the retardation effects in the two-body collision integral $J_c(\{\delta f\})$. We point out also that the relaxation time $\tau_c(\omega_0,T)$ is different for the isovector and isoscalar excitations. The components of the distorted distribution function δf in momentum space with multipolarities $\ell=0$ and $\ell=1$ do not contribute to the collision integral $J_c(\{\delta f\})$ for isoscalar excitation because of the above-mentioned conservation of the number of particles and the total momentum. However, in the case of the isovector mode there is the motion of protons against neutrons in phase space without the violation of the conservation of the total momentum and the dipole $\ell=1$ distortion of the distribution function $\delta f^{(-)}$ also gives nonzero contribution to the collision integral $J^{(-)}(\{\delta f_p,\delta f_n\})$.

To begin with, we will consider in Eq. (1) the Fermi-surface distortion with multipolarities $\ell \leqslant 2$, assuming the following expression for $\delta f(\mathbf{r},\mathbf{p},t)$:

$$\delta f(\mathbf{r},\mathbf{p},t) = \sum_{\ell=0}^{2} \delta f_{\ell}(\mathbf{r},\mathbf{p},t) \equiv \sum_{\ell=0}^{2} \delta \tilde{f}_{\ell}(\mathbf{r},p,t) Y_{\ell 0}(\hat{p}). \tag{9}$$

Here, $Y_{\ell m}(\hat{p})$ is the spherical harmonic function and $\hat{p}=\mathbf{p}/p$. In accordance with the result of our earlier work [6], the widths of slow damped collective vibrations are determined by the relaxation tensor $Q_{\alpha\beta}$ given by

$$Q_{\alpha\beta} = \frac{g}{m}\int \frac{d\mathbf{p}}{(2\pi\hbar)^3}(p_\alpha - m u_\alpha)(p_\beta - m u_\beta) J(\{\delta f\}). \tag{10}$$

We introduce the dynamical component of the pressure tensor $P'_{\alpha\beta}$ associated with dissipative processes and given by

$$P'_{\alpha\beta} = \frac{g}{m}\int \frac{d\mathbf{p}}{(2\pi\hbar)^3}(p_\alpha - m u_\alpha)(p_\beta - m u_\beta)\,\delta f. \tag{11}$$

The quantity u_α is the Cartesian component of the velocity field $\mathbf{u}$ and ρ is the particle density.

Substituting Eqs. (6)–(8) into Eq. (10) and evaluating the contribution from J_c using the method discussed in Ref. [6], one obtains

$$Q_{\alpha\beta} = -P'_{\alpha\beta}/\tau_2(\omega_0,T) - P'_{\alpha\beta}/\tau_s - P'_{\alpha\beta}/\tau_{\uparrow}. \tag{12}$$

Thus,

$$Q_{\alpha\beta} = -P'_{\alpha\beta}/\tau_{\mathrm{eff}}, \tag{13}$$

where τ_{eff} is the effective relaxation time

$$1/\tau_{\mathrm{eff}} = 1/\tau_2(\omega_0,T) + 1/\tau_s + 1/\tau_{\uparrow}. \tag{14}$$

Here, $\tau_2(\omega_0,T)$ is the two-body collision relaxation time for the collective vibration with eigenfrequency ω_0 and at temperature T in the case of quadrupole distortion of the Fermi surface. It has the following form [6,11,30–34]

$$1/\tau_2(\omega_0,T) = [1 + C_\omega(\hbar\omega_0/2\pi T)^2]/\tilde{\tau}_2(T), \tag{15}$$

which is valid at $T,\hbar\omega_0 \ll \mu \approx \epsilon_F$. The factor C_ω in Eq. (15) defines the value of the relaxation time τ_2 in the quantum region $\hbar\omega_0 \gg T$. The magnitude of this factor is discussed in [6] and it is equal to

$$C_\omega = 1, \tag{16}$$

if one follows Landau's prescription [32]. The relaxation time $\tau_2(\omega_0,T)$ is frequency and temperature dependent. The frequency dependence of $\tau_2(\omega_0,T)$ is due to the memory effect in the collision integral (8). The temperature dependence in Eq. (15) arises from the smeared out behavior of the equilibrium distribution function f_0, see Eq. (34), near the Fermi momentum. The quantity $\tilde{\tau}_2(T)$ in Eq. (15) is the thermal relaxation time that will be discussed later (see Sec. III).

Equations (1), (9), and (10) coincide with the analogous Eqs. (1), (2), and (21) of Ref. [6], replacing $J(\{\delta f\})$ by $St(t)$. Therefore, the expression for the intrinsic GMR width Γ has to have the form given by Eq. (38) of Ref. [6] by replacing τ_2 with τ_{eff}. Thus,

$$\Gamma = 2\,\tilde{c}B(S_r)\hbar\omega_0 \frac{\omega_0\tau_{\mathrm{eff}}}{1+\tilde{c}(\omega_0\tau_{\mathrm{eff}})^2}, \quad \tilde{c} \equiv B(S_f)/B(S_r). \tag{17}$$

Here the function $B(S)$ is given by

$$B(S)^{-1} \equiv 15S^2/2 = B(S_r)^{-1} + [B(S_f)^{-1} - B(S_r)^{-1}]/[1+(\omega_0\tau_{\mathrm{eff}})^2]. \tag{18}$$

The quantity $S \equiv \omega_0/v_F k$ is the dimensionless velocity of sound wave in the Fermi liquid and $v_F = p_F/m^*$ is the Fermi velocity. In the case of the quadrupole distortion of the Fermi surface, see Eq. (9), one has [6,35]

$$S^2 = S_r^2 + [S_f^2 - S_r^2]/[1+(\omega_0\tau_{\mathrm{eff}})^2], \tag{19}$$

with

$$S_f \equiv \sqrt{K/9mv_F^2}, \quad S_r \equiv \sqrt{[K + 8\,(\epsilon_{\mathrm{kin}}/\rho)_{\mathrm{eq}}]/9mv_F^2}. \tag{20}$$

Here K is the adiabatic incompressibility modulus given by Eq. (4), the subscript eq means that the corresponding values are taken at equilibrium, and ϵ_{kin} is the kinetic energy density in the case of a spherical Fermi surface. The quantities S_f and S_r are the velocities of the first and the zero sounds (in units of v_F and under the condition of quadrupole distortion of the Fermi surface), respectively. It is necessary to emphasize that, as it follows from Eqs. (19) and (20), a condition for propagation of the zero and the first sounds in a finite system is determined by the magnitude of the parameter $\omega_0\tau_{\rm eff}$. The effective relaxation time $\tau_{\rm eff}$, see Eq. (14), includes the contributions from the two-body collisions, the one-body wall dissipation, and the particle emission. In contrast, in an infinite Fermi liquid the analogous condition for the propagation of sound waves is governed by the two-body collision only [36,37]. We point out also that expression (17) for the intrinsic width of GMR is obtained within the many-particle model where the infinite-matter solutions to the Landau-Vlasov equation are completed with some boundary conditions at a sharp edge [6,30,31].

As mentioned above [see Eq. (9)], Eq. (17) is valid in the case of the quadrupole distortion of the Fermi surface. In the same manner as in Ref. [6], the approximate formula (17) for the width can be deduced by taking into account other multipolarities of the dynamical distortion of the Fermi surface. Using Eqs. (45), (51), and (52) from Ref. [6] one can write

$$\Gamma=2\,qE\,\frac{E\tau/\hbar}{1+q(E\tau/\hbar)^2},\quad E\equiv\hbar\omega_0, \tag{21}$$

where E is the energy of the GMR. The cutoff factor q and the relaxation time τ are given by

$$q\approx 1/6S_f^2\approx 1/2(1+F_0)(1+F_1/3), \tag{22}$$

$$1/\tau=1/\tau_c+1/\tau_s+1/\tau_\uparrow, \tag{23}$$

with [see also Eq. (15)]

$$\tau_c\equiv d_2\tau_2(\omega_0,T)=d_2\tilde{\tau}_2(T)/[1+C_\omega(E/2\pi T)^2] \\ \equiv\tilde{\tau}_c(T)/[1+C_\omega(E/2\pi T)^2]. \tag{24}$$

Here $\tilde{\tau}_c(T)\equiv d_2\tilde{\tau}_2(T)$ and the quantity d_2 is

$$d_2=\langle w\Phi_2\rangle/\langle w\rangle, \tag{25}$$

where w is the probability for scattering of two indistinguishable particles near the Fermi surface [36]. The function

$$\Phi_2(\theta,\varphi)=3\sin^4\frac{\theta}{2}\sin^2\varphi \tag{26}$$

defines the angular constraints for nucleon scattering within the distorted layer of the Fermi surface with $\ell=2$ and the brackets denote the averaging of the form

$$\langle\Psi\rangle\equiv\frac{1}{2\pi}\int_0^\pi d\theta\int_0^\pi d\varphi\,\frac{\sin\theta}{\cos(\theta/2)}\Psi(\theta,\varphi) \tag{27}$$

over angles θ and φ which are determined by the momentum $\mathbf{p}_j$ $(j=1\text{–}4)$ of the colliding particles. Namely,

$$\cos\theta=(\hat{p}_1\cdot\hat{p}_2), \\ \cos\varphi=[\hat{p}_1\times\hat{p}_2]\cdot[\hat{p}_3\times\hat{p}_4]/|[\hat{p}_1\times\hat{p}_2]||[\hat{p}_3\times\hat{p}_4]|. \tag{28}$$

As can be found from Eqs. (25)–(27), if the probability w of the two-particle scattering is isotropic in space, then the magnitude of d_2 is

$$d_2=4/5,$$

and everywhere below we will use this value.

It follows from Eqs. (21) and (23) that in the long relaxation time regime, i.e., at $E\tau/\hbar\gg 1$, the expression for intrinsic width Γ has the additive form

$$\Gamma\equiv\Gamma_0=2\hbar/\tau=\Gamma_c+\Gamma_s+\Gamma_\uparrow, \tag{29}$$

where

$$\Gamma_c\equiv 2\hbar/\tau_c,\quad \Gamma_s\equiv 2\hbar/\tau_s,\quad \Gamma_\uparrow\equiv 2\hbar/\tau_\uparrow. \tag{30}$$

The expressions (29) and (30) for the width correspond to those in the relaxation rate approximation when the interrelationship between different dissipation channels is ignored.

As can be seen from Eqs. (21) and (23), in a general case, the total intrinsic width is not represented as a sum of the partial widths of separate channels in spite of the fact that relation (6) was used. This is a peculiarity of the collisional Landau-Vlasov equation where the Fermi-surface distortion effect influences both the self-consistent mean field and the memory effect in the relaxation processes.

III. DAMPING PROPERTIES OF THE GIANT MULTIPOLE RESONANCES

The values of the GMR energy E, the relaxation times in different damping channels τ_c, τ_s, and $\tau_\uparrow$ and the cutoff factor q are required for calculations of the intrinsic width Γ. As the GMR energy E in Eqs. (21) and (24), we use the phenomenological A dependence of E obtained from a fit to the experimental data [38–41]. The value of q is determined by the Landau parameters F_0 and F_1 of the nucleon-nucleon interaction, see Eq. (22). As mentioned earlier, to apply Eq. (21) for a description of both the isoscalar and the isovector GMR one has to use two sets of the Landau parameters. Namely, $F_0^{(+)}$ and $F_1^{(+)}$ for the isoscalar GMR and $F_0^{(-)}$ and $F_1^{(-)}$ for the isovector GMR, see Eq. (3). We will use for $F_0^{(\pm)}$ and $F_1^{(\pm)}$ the values determined for nuclear matter [42,43].

(*i*) *Collisional relaxation time.* The collisional relaxation time τ_c, Eq. (24), depends on the thermal relaxation time $\tilde{\tau}_2(T)$. In a homogeneous Fermi system the relaxation time $\tilde{\tau}_2(T)$ is given by [30,31,44,45]

$$\frac{1}{\tilde{\tau}_2(T)}\equiv\int d\mathbf{p}J_c^{UUN}(\{\delta f\})Y_{20}(\hat{p})\Bigg/\int d\mathbf{p}\,\delta f Y_{20}(\hat{p}). \tag{31}$$

Here, $J_c^{UUN}(\{\delta f\})$ is the linearized Uehling-Uhlenbeck-Nordheim collision integral in the absence of the retardation effects [37]:

$$J_c^{UUN}(\{\delta f\}) \equiv \lim_{T \gg \hbar\omega_0} J_c(\{\delta f\}). \quad (32)$$

The dynamical component δf of the distribution function has the form

$$\delta f(\mathbf{p},\mathbf{r},t) = -(\partial f_0/\partial\epsilon)\psi(\hat{p},\mathbf{r})\exp(i\omega t). \quad (33)$$

Here the function ψ depends only on the direction of the momentum, and the equilibrium component f_0 of the Wigner distribution function is taken as the Fermi distribution

$$f_0 = 1/(1+\exp\{(\epsilon-\mu)/T\}), \quad (34)$$

depending on the quasiparticle energy $\epsilon = p^2/(2m^*)$.

The magnitude of $\tilde{\tau}_2(T)$ of Eq. (31) can be represented in the following general form at low temperatures $T \ll \mu$ [6]:

$$\tilde{\tau}_2(T)/\hbar = \alpha T^{-2} \quad (T,\alpha \text{ in MeV}), \quad (35)$$

where the quantity α in the case of the isoscalar mode is given by

$$\alpha^{(+)} = 5\epsilon_F^2/4\pi^2\hbar\rho v_F\sigma_{\rm av}. \quad (36)$$

It is determined by the in-media spin-isospin averaged nucleon-nucleon cross section $\sigma_{\rm av}$

$$\sigma_{\rm av} = (\sigma_{\rm pp} + \sigma_{\rm nn} + 2\,\sigma_{\rm pn})/4.$$

Here, $\sigma_{\rm pp}, \sigma_{\rm nn}$, and $\sigma_{\rm pn}$ are the in-media cross sections for nucleon pairs with relative kinetic energy close to the Fermi energy. In Eq. (36) ρ is the nuclear matter density. In the case of the isovector mode, there is an additional contribution from the dipole distortion of the Fermi surface and the quantity α is given by

$$\alpha^{(-)} = 5\,\epsilon_F^2/4\pi^2\hbar\rho v_F(\sigma_{\rm av} + 5\,\sigma^{(-)}/3), \quad (37)$$

where $\sigma^{(-)} = \sigma_{\rm pn}/2$. Thus, in the case of the dipole isovector mode there is an enhancement of the collisional width due to nonconservation of the isovector current in the neutron-proton collisions [34].

The assessments of the relaxation time $\tilde{\tau}_2(T)$, and thereby α, using the free space nucleon-nucleon cross sections [11,45–48] $\sigma_{pp} = \sigma_{nn} = 2.5$–$2.7$ fm^2 and $\sigma_{pn} = \sigma_{np} = 4.8$–$5.0$ fm^2 and the nuclear matter density $\rho = 0.17$ fm^{-3} give

$$\alpha^{(+)} = 4.6 - 4.9 \text{ MeV and } \alpha^{(-)} = 2.2 - 2.3 \text{ MeV}. \quad (38)$$

Due to the Pauli blocking effect it is expected that the collision probability w in nuclear matter should be lower than the one in free space. We will follow Refs. [47–51] and use the value of the nucleon-nucleon cross section in-medium to be smaller than the cross section in free space by a factor of about of 2. Thus, we will use

$$\alpha^{(+)} = 9.2 \text{ MeV and } \alpha^{(-)} = 4.6 \text{ MeV} \quad (39)$$

as the more realistic values of α.

An independent assessment of the thermal relaxation time $\tilde{\tau}_2(T)$ can be obtained by using the lifetime $\tau_p(\epsilon,T)$ of the single particle excited state with energy ϵ in a Fermi system having temperature T [37,52,53]. Within the framework of the Fermi-liquid transport theory [37,52,53], the lifetime $\tau_p(\epsilon,T)$ is related to the linearized Uehling-Uhlenbeck-Nordheim collision integral (32) as

$$\frac{1}{\tau_p(\epsilon,T)} \equiv J_c^{UUN}(\{\delta f\})/\delta f. \quad (40)$$

Using Eqs. (31)–(34) and (40) and employing a standard transformation of the Fermi-liquid theory [36] for the calculation of the integrals over momenta in Eq. (31), we obtain

$$1/\tilde{\tau}_2(T) = -d_2\int_0^\infty d\epsilon(\partial f_0/\partial\epsilon)[1/\tau_p(\epsilon,T)], \quad (41)$$

or

$$1/\tilde{\tau}_c(T) = -\int_0^\infty d\epsilon(\partial f_0/\partial\epsilon)[1/\tau_p(\epsilon,T)]. \quad (42)$$

Notice, as it follows from the definition of $\tau_p(\epsilon,T)$, Eq. (40), the relaxation time $\tau_p(\epsilon,T)$ in the low temperature regime $T \ll \mu$ is given by (see, for example, Refs. [37,54])

$$\frac{1}{\tau_p(\epsilon,T)} = \beta\{(\epsilon-\mu)^2 + \pi T^2\}, \quad (43)$$

where the parameter β does not depend on T and ϵ. Substituting Eqs. (43) and (34) into Eq. (41) and comparing the result of the integration with the expression (35), we find the following relationship between the parameters α and β

$$\alpha = 3/4\pi^2\hbar d_2\beta \approx 15/16\pi^2\hbar\beta. \quad (44)$$

We will now obtain an assessment of the parameter β using the exciton model of nuclear reactions [55]. Following Refs. [56,57], the lifetime $\tau_p(\epsilon,T=0)$ at zero temperature can be identified with the lifetime $\tau^{(1)}(U)$ of the one-exciton state having one particle and no hole at the excitation energy $U = \epsilon - \epsilon_F$,

$$\tau_p(\epsilon,T=0) = \tau^{(1)}(U) \equiv 1/\lambda_+^{(1)}(U = \epsilon - \epsilon_F). \quad (45)$$

Here,

$$\lambda_+^{(n)}(U) \equiv \lambda_+(n_p p, n_h h; U) = \frac{2\pi}{\hbar}|M|^2\Delta\rho_f[(n_p+1)p,(n_h+1)h;U] \quad (46)$$

is the rate for the particle-hole pair creation due to the two-body collisions starting from the initial configuration containing n_p particles and n_h holes, i.e., $n = n_p + n_h$ excitons. The quantity $|M|^2$ in Eq. (46) is the mean square matrix element for the residual two-body interactions and $\Delta\rho_f[(n_p+1)p,(n_h+1)h;U]$ is the density of the available final states for a transition $n \to n+2$. In the equidistant-spacing model for the single particle level density $g_{\rm s.p.}$ one has [56–58]

$$\Delta\rho_f[(n_p+1)p,(n_h+1)h;U] = \frac{g_{\rm s.p.}}{2}\frac{(g_{\rm s.p.}U)^2}{n_p+n_h+1}. \quad (47)$$

A comparison of Eq. (43) and Eqs. (45)–(47) leads to

TABLE I. Theoretical assessments for parameters ξ^L for isoscalar $L^\pi=0^+$, 2^+, and 3^- and isovector $L^\pi=1^-$ GMR at given multipolarity L and parity π.

L^π	0^+	1^-	2^+	3^-	Theory of one-body dissipation
$\xi^L_{(1)}$		1.00	0.50	0.33	wall, Refs. [60–62]
$\xi^L_{(2)}$			6.00	1.00	modified wall; $A=125$, Refs. [60,63]
$\xi^L_{(3)}$			1.85	1.24	rescaling of wall, Refs. [64,65]
$\xi^L_{(4)}$			5.00	3.33	RPA wall, Refs. [7–9]
$\xi^L_{(5)}$	5.60	1.81	2.43	1.00	fit; $\alpha=9.2$ MeV, $q^{(+)}=0.282$, $q^{(-)}=0.192$
$\xi^L_{(6)}$	<0	3.20	4.02	2.09	fit; $\alpha=4.6$ MeV, $q^{(+)}=0.282$, $q^{(-)}=0.192$

$$\beta=\frac{\pi}{2\hbar}|M|^2 g^3_{\text{s.p.}}. \tag{48}$$

Inserting this expression into Eq. (44), we find the relationship between the parameter α and the mean square matrix element $|M|^2$

$$\alpha=3/2\pi^3 d_2 g^3_{\text{s.p.}}|M|^2\approx 15/(2\pi g_{\text{s.p.}})^3|M|^2. \tag{49}$$

Substituting $g_{\text{s.p.}}=(6/\pi^2)(A/7.5)$ MeV^{-1} and $|M|^2=(15.2/A^3)$ MeV^2 (see Refs. [56,57,59]) in Eq. (49), we find

$$\alpha\approx 7.5 \text{ MeV}, \tag{50}$$

which is in a good agreement with that given by Eq. (39).

(*ii*) *One-body relaxation time.* We will now discuss a choice of the relaxation time $\tau_s\equiv\tau_s^L$ which determines the contribution of the one-body wall dissipation to the damping of the GMR with multipolarity L. The relaxation time τ_s^L is related to the partial width Γ_s^L of the GMR arising from the one-body dissipation by

$$\tau_s^L=2\hbar/\Gamma_s^L. \tag{51}$$

A number of authors estimated earlier the one-body GMR width Γ_s^L or the associated friction coefficient γ_s^L within both the classical and the quantum approaches [7–9,20,60–67]. It was shown that Γ_s^L is proportional to the relative weight of the surface region in the system. We will represent below the one-body relaxation time τ_s^L in the form

$$\tau_s^L=\xi^L\tau_d, \quad \tau_d\equiv 2R/\overline{v}, \tag{52}$$

where $R_0=r_0A^{1/3}$ is the nuclear radius, τ_d is the time of flight of the free nucleon through the nuclear diameter. The factor ξ^L in Eq. (52) specifies the difference between τ_d and τ_s^L and depends on the model of the one-body dissipation. The quantity $\overline{v}$ is the average velocity of the nucleon incorporating also the temperature effects. It is given by [9]

$$\overline{v}=(3v_F/4)[1+(\pi^2/6)(T/\epsilon_F)^2]. \tag{53}$$

In Table I we show a list of the parameters ξ^L which were derived through Eqs. (51) and (52) using different estimates for Γ_s^L.

The coefficients $\xi^L_{(1)}$ were estimated by means of the expression for the one-body width from Refs. [60,61] where the width Γ_s^L was calculated employing the classical Swiatecki formula [62] under the assumption of incompressible, irrotational nuclear flow. As a result, the factor $\xi^L_{(1)}$ has the following closed form

$$\xi^L_{(1)}=1/L, \quad L\geqslant 1. \tag{54}$$

The values $\xi^L_{(2)}$ have been found by using the expression for Γ_s^L from Ref. [60]. The isoscalar resonance width Γ_s^L was here evaluated as pointed to above but with the modified expression (from Ref. [63]) for the friction coefficients $\gamma_s^L=\Gamma_s^L B_L/\hbar$, where B_L is the mass parameter for small shape oscillations in the liquid drop model [68]. The modified friction coefficient γ_s^L takes into account a consistent description of the nuclear surface motion and the internal motion of particles. The corresponding coefficient $\xi^L_{(2)}$ is given by

$$\xi^L_{(2)}=\xi^L_{(1)}(R/\lambda)^2/(L-1)^2=\left(\frac{R}{\lambda}\right)^2\frac{1}{L(L-1)^2}, \quad L\geqslant 2, \tag{55}$$

where $\lambda=1.7\pm0.3$ fm [60] is the effective distance specifying the magnitude of the dissipation.

The coefficients $\xi^L_{(3)}$ are extracted through the use of the expressions for the one-body widths from Ref. [64]. These values of Γ_s^L are less than those for wall dissipation of Eq. (54) by the factor $k_s=0.27$. The magnitude of k_s was found from a fit to the experimental data for the widths of the isoscalar giant quadrupole and octupole resonances. Thus,

$$\xi^L_{(3)}=\xi^L_{(1)}/k_s\approx 3.7/L. \tag{56}$$

The values $\xi^L_{(4)}$ were estimated by means of Γ_s^L from Ref. [7] where the one-body width of the GMR was identified with its fragmentation width in the random phase approximation. The quantal calculation of Γ_s^L in a simplified version of the RPA [7–9] shows a significant enhancement of the ξ^L (or the corresponding reduction of the width Γ_s^L) in comparison with the ''wall'' value given by Eq. (54). The following extension of the one-body width Γ_s^L can be derived within the RPA (see Refs. [7–9]):

$$\Gamma_s^L=\frac{L}{R}G(E/\epsilon_F,B/\epsilon_F)\overline{v}, \tag{57}$$

where B is the nucleon binding energy. The function $G(E/\epsilon_F,B/\epsilon_F)$ has a threshold behavior [$G(E/\epsilon_F,B/\epsilon_F)=0$ at $E<B$] and depends on the reflection

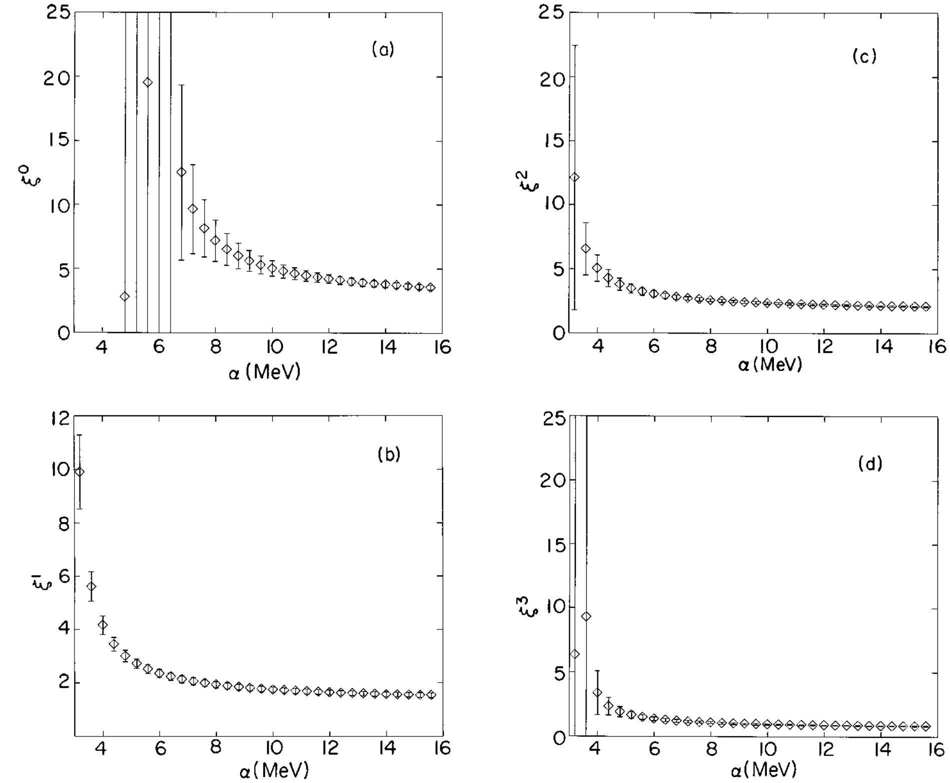

FIG. 1. The dependence of the factors ξ^L, see Eq. (52), for the one-body relaxation time on the parameter α of the two-body thermal relaxation time, Eq. (36). The correspondence between the figures and giant resonances is the following: $1a-0^+$, $1b-1^-$, $1c-2^+$, and $1d-3^-$. The following values of the Landau amplitudes were used: $F_0^{(-)}=1.6$ and $F_1^{(-)}=0$ [42], i.e., $q^{(-)}=0.192(S_f^{(-)}=0.931)$ for the isovector GDR (1^-) and $F_0^{(+)}=0.77$ and $F_1^{(+)}=0$ [43], i.e., $q^{(+)}=0.283(S_f^{(+)}=0.768)$ for the isoscalar GMR $(0^+,2^+,3^-)$ were used.

conditions of particles on the potential wall. Only in the case of the fully-reflecting potential wall and in the adiabatic limit $E/\epsilon_F\rightarrow 0$ and at $B/\epsilon_F\rightarrow 0$, one has $G(E/\epsilon_F,B/\epsilon_F)=1$ and the quantal RPA expression (57) coincides with the classical wall formula [20]. Numerical calculation shows [7] that the one-body widths Γ_s^L are about 10 times smaller than those obtained within the wall model, Eq. (54). Thus, we will use for $\xi_{(4)}^L$ the following estimate

$$\xi_{(4)}^L\approx 10/L. \tag{58}$$

The values $\xi_{(5)}^L$ and $\xi_{(6)}^L$ in Table I were obtained by a fit to experimental data, as described in the next section.

(iii) *Relaxation time due to particle emission.* We will estimate the particle emission relaxation time $\tau_\uparrow$ taking into account the thermal particle evaporation only. The particle emission relaxation time $\tau_\uparrow$ is connected to the corresponding width $\Gamma_\uparrow$ by the usual relationship $\tau_\uparrow=2\hbar/\Gamma_\uparrow$. We use for $\Gamma_\uparrow$ the evaporation formula for neutrons [4]

$$\Gamma_\uparrow\approx\Gamma_n=\frac{2mR^2}{\pi\hbar^2}T^2\exp(-B_n/T), \tag{59}$$

where B_n is the neutron binding energy.

IV. NUMERICAL RESULTS AND DISCUSSION

In heavy and medium nuclei the contribution of particle emission from the GMR to the width is small enough at low temperature [4] and can be neglected. We will therefore take into account only the one- and two-body channels of damping in Eqs. (21) and (29) at $T=0$.

Figure 1 shows the dependence of the factors ξ^L for the one-body relaxation times τ_s^L of Eq. (52) on the parameter α in the collisional relaxation time τ_c, see Eqs. (24) and (35). This dependence was obtained from a fit of Eq. (21) to the experimental data [38–41] for the GMR widths in cold nuclei with mass numbers $50\leqslant A\leqslant 260$. As can be seen from these figures, for all GMR (excluding the monopole) the quantities ξ^L are almost independent of α when the condition $\alpha\geqslant 6.5$ MeV is fulfilled. We point out that we do not distinguish here between the isoscalar and isovector modes because the quantity α is the adjusted parameter, i.e., the relations (36) and (37) are not used for the evaluation of α.

In Fig. 2 we show the intrinsic GMR widths and the one- and the two-body contributions to them at $T=0$ as functions of mass number. The total intrinsic widths are found by means of Eq. (21) with $1/\tau=1/\tau_c+1/\tau_s$. The different partial contributions to the GMR widths are also evaluated by Eq. (21), but with the relaxation times corresponding to the selected partial damping channel. We used for α the values of 9.2 and 4.6 MeV and the corresponding values of $\xi_{(5)}^L$ and $\xi_{(6)}^L$ were taken from Table I. These values of $\xi_{(5)}^L$ and $\xi_{(6)}^L$ were deduced from Fig. 1. As can be seen from Fig. 2, the contribution of collisional damping (curves 2 and 2′) to the GMR widths does not exceed ~60% of the experimental values for $\alpha\geqslant 4.6$ MeV, irrespective of the type of the GMR and the mass number A. This contribution decreases with increasing α.

In Fig. 3 we show the intrinsic width of the giant dipole resonance (GDR) in the nucleus ^{112}Sn. The experimental data were taken from Ref. [3]. Considering the experimental data we assumed that the energy E of the GDR is independent of temperature and equals 15.6 MeV. We used the relation $U=aT^2$ between the temperature T and the excitation energy U, deduced in the Fermi system, where a is the level density parameter. We adopted the value $a=A/8$ MeV^{-1}.

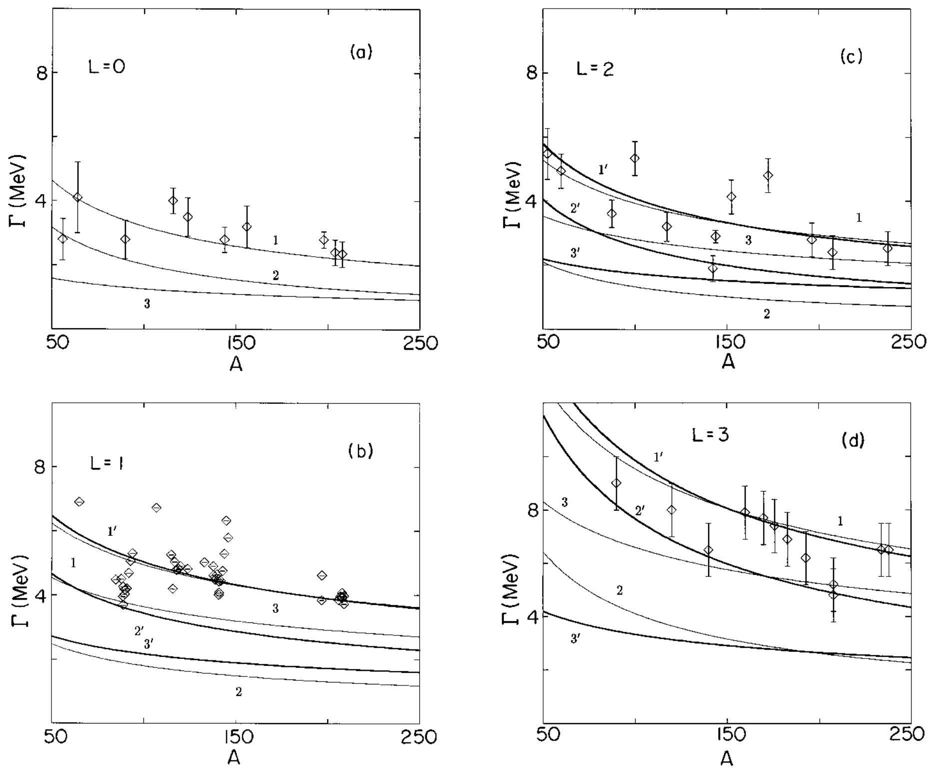

FIG. 2. The intrinsic GMR widths and the corresponding one- and the two-body contributions at $T=0$ as functions of mass number. The correspondence between figures and giant resonances is the same as in Fig. 1. Curves 1 and $1'$ correspond to the total width. Curves 2, $2'$ and 3, $3'$ correspond to the two- and the one-body contributions, respectively. The value $\alpha=9.2$ MeV is used for curves 1, 2, and 3 and $\alpha=4.6$ MeV for curves $1'$, $2'$, and $3'$. The values of ξ^L corresponding to given values of α are given in Table I as $\xi^L_{(5)}$ and $\xi^L_{(6)}$. The experimental data were taken from Refs. [38–41].

We have also taken into account the contribution to the total intrinsic widths from particle emission, see Eq. (59), with $B_n=10.8$ MeV for ^{112}Sn nucleus. As is seen from Fig. 3, the expression (21) (curves 1 and 5) leads to a smoother behavior of the total intrinsic width with increasing excitation energy as compared with the prediction of the zero sound model given by Eqs. (29) and (30) (curve 6). Our calculation of the intrinsic width Γ for the isovector GDR in hot nucleus ^{112}Sn confirms a saturation effect in the energy dependence of Γ. However, we observe a systematic deviation of the evaluated width with respect to the experimental data. This deviation can be reduced by varying the Landau parameter $F_1^{(-)}$, and thereby the cutoff factor q in Eq. (21). We recall that in Ref. [6] we only considered the contribution of the collisional damping to $\Gamma(U)$ and showed that an increase in the cutoff parameter $q^{(-)}$ improves the agreement with experiment (see Fig. 3 of Ref. [6]).

In Fig. 4 the parameter $\omega_0\tau$ is plotted as a function of excitation energy for the giant dipole resonance in ^{112}Sn nucleus. We notice that the parameter $\omega_0\tau$ governed the transition from the zero sound regime to the first sound regime. Due to the presence of the one-body damping and the particle emission in a finite system, the value of $\omega_0\tau$ is different from its value $\omega_0\tau_c$ in an infinite nuclear matter. This can be seen by comparing curves 1 and 2 with 3 in Fig. 4. In principle, the situation may occur when in an infinite nuclear matter $\omega_0\tau\equiv\omega_0\tau_c\approx1$ but in a finite nucleus the short relaxation time regime is realized with $\omega_0\tau\ll1$. In the short re-

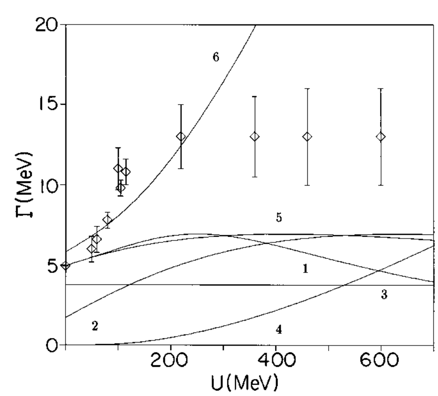

FIG. 3. The intrinsic width of the giant dipole resonance in the nucleus Sn as a function of excitation energy. The experimental data were taken from Ref. [3]. Curve 1 corresponds to the total width. Curves 2, 3, and 4 correspond to the two- , one-body, and particle emission ($\Gamma_\uparrow\approx\Gamma_{\text{neutron}}$) contributions, respectively. Curve 5 is the width without the contribution of particle emission. Curve 6 is the width in the zero sound approximation of Eq. (29). The values of $\alpha=4.6$MeV and $\xi^1=3.2$ were used. This value of ξ^1 is the result of a fit to experimental data at $T=0$.

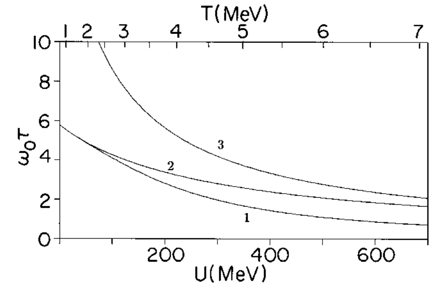

FIG. 4. The excitation energy dependence of $(\omega_0\tau)$ for the giant dipole resonance in ^{112}Sn nucleus. Curve 1 corresponds to the total effective relaxation time. Curve 2 is obtained without taking into account particle emission. Curve 3 corresponds to $(\omega_0\tau)$ with two-body relaxation time as τ.

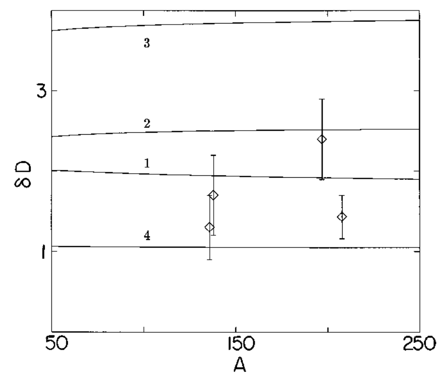

FIG. 5. Atomic mass number dependence of the ratio $\delta D = \Gamma_{\rm double}/\Gamma$. Curves 1 and 2 correspond to the contributions of all processes to the relaxation time. Curves 3 and 4 correspond to the calculations with the two- and the one-body relaxation time, respectively. The values of $\alpha = 9.2$ MeV and $\xi^1 = 1.81$ were used for curves 1, 3, and 4 and $\alpha = 4.6$ MeV and $\xi^1 = 3.2$ were used for curve 2. The experimental date were taken from Ref. [69].

laxation time regime (the first sound regime), the expression for the width has the form

$$\Gamma \equiv \Gamma_f = 2qE^2\tau/\hbar, \qquad (60)$$

as it follows from Eq. (21). This expression is not represented as a sum of the partial widths connected with different damping channels as it is in the long relaxation time regime, see Eq. (29). Notice that allowing for the particle emission leads to a reduction of the total intrinsic width with increasing temperature. In particular, the behavior of curve 1 in Fig. 3 for excitation energies $U \geqslant 300$ MeV is connected to the fact that $\omega_0\tau \approx 1$ in this energy range, see Fig. 4. In this case, the regime of the collective motion is more similar to the short relaxation time regime with the width given by Eq. (60). Since the condition $\tau \approx \tau_\uparrow = 2\hbar/\Gamma_\uparrow$ is fulfilled at $U \geqslant 300$ MeV and the magnitude of $\Gamma_\uparrow$ increases with U, the values of the relaxation time τ and the width Γ will decrease because of Eq. (60).

We will now discuss the intrinsic width $\Gamma_{\rm double}$ of the double isovector giant dipole resonance (DGDR) which follows from our model. In Fig. 5 we show the mass number dependence of the ratio $\delta D = \Gamma_{\rm double}/\Gamma$. Both values of $\Gamma_{\rm double}$ and Γ were obtained from Eq. (21) with $E = 2\hbar\omega_{1^-}$ and $E = \hbar\omega_{1^-}$, respectively, where $\hbar\omega_{1^-}$ is the energy of the isovector GDR. As can be seen from Fig. 5, the magnitude of δD is rather sensitive to the value of α which characterizes the contribution from the two-body collisions to the intrinsic width. One has $\delta D \approx 4$ when the two-body damping dominates and $\delta D \approx 1$ if only the one-body channel is taken into account.

Finally, we want to note that the total intrinsic width given by Eq. (21) has a bell-shaped form as a function of $x \equiv E\tau/\hbar$. The width Γ is peaked at $x \equiv x_0 = q^{-1/2}$ and the maximum value of Γ is $\Gamma_{\rm max} = Eq^{1/2}$. It is easy to see that x_0 represents the crossing point of both curves $\Gamma_0(x)$ and $\Gamma_f(x)$ given by Eqs. (29) and (60). Due to this fact the condition

$$\omega_0\tau \equiv x_0 = (\omega_0\tau)_0 = q^{-1/2} \qquad (61)$$

can be used as the condition for the transition from the long- to the short-relaxation time regimes. The magnitude of the intrinsic width Γ decreases when the parameter $\omega_0\tau$ exceeds $(\omega_0\tau)_0$. We have $(\omega_0\tau)_0 = 2.28$ for the isovector GDR at the realistic value of $q^{(-)} = 0.192$. As it follows from Fig. 4, such a value of $\omega_0\tau$ can be reached at temperature $T \equiv T_{\rm tr} \approx 4.5$ MeV. If the relation (61) with $q = 1$ and $\tau = \tau_c$ is used as the condition for the transition between different regimes of propagation of sound wave, see Ref. [37], then the transition between both regimes occurs at higher temperature $T_{\rm tr} \approx 10$ MeV.

V. CONCLUSION AND SUMMARY

We have performed a phenomenological analysis of the interplay between different damping mechanisms of collective motion in hot nuclei. Three main channels of dissipation have been taken into account: the two-body and one-body collision channels and the particle emission. The two-body collision channel describes the relaxation on the deformed Fermi surface due to the interparticle collisions. The origin of the one-body dissipation is the collision of the nucleons with a moving nuclear surface. The nuclear surface is considered here as the collective degree of freedom. We pointed out that the one-body dissipation gives a macroscopic description of the fragmentation width of the GMR appearing in a microscopic approach like the RPA. The particle emission channel accounts for the particle evaporation. Its contribution to the total GMR width grows very fast with the excitation energy, see Eq. (59), and becomes comparable to both other channels at temperature about 6 MeV, see Fig. 3.

We have used the approximation of independent dissipation rates, Eq. (6), for all mentioned dissipation channels. In spite of this fact, the intrinsic width Γ of GMR, Eq. (21), can not be represented in an additive form with respect to each channel. This is due to the fact that the Fermi-surface distortion effect influences both the self-consistent mean field and the memory effect at the relaxation processes. The additive form of the intrinsic width, see Eq. (29), is achieved in the limit of the long relaxation time $\omega\tau \gg 1$.

The contribution of the two-body dissipation channel into the intrinsic width Γ is determined by the parameter α in the collisional relaxation time $\tilde{\tau}_2$, Eq. (35), and depends on the in-media nucleon-nucleon cross section, see Eqs. (36) and (37). The different assessments for the magnitude of α, within the framework of the kinetic theory of nuclear Fermi liquid, give [11,45,46,70,71] 2.4 MeV $< \alpha <$ 19.2 MeV. Our independent estimate for α from the exciton model of nuclear reactions leads to $\alpha = 7.5$ MeV. An important ingredient of our consideration is the collisional memory effect which is manifested in the dependence of the collisional width Γ_c on the frequency of collective vibrations. Due to this fact, our expression (21) for the intrinsic width is valid not only in the regime of rare collisions but also in the transition region from the zero-sound to the first-sound (hydrodynamic) regime. The latest is important in heated nuclei where the fast collision regime can be achieved.

We derived the relative contribution of the one-body dissipation channel [parameter ξ^L in Eq. (52)] from a fit of the intrinsic width Γ to the experimental data in cold nuclei. Our calculations show a weak sensitivity of ξ^L to the choice of the parameter α at $\alpha \geqslant 6.5$ MeV. In the case of the isoscalar GMR the parameter $\xi_{(5)}^{L}$ obtained for a realistic $\alpha = 9.2$ MeV (see Table I) differs significantly from the predictions of the wall formula as well as its modifications [7–9,60–64].

The retardation effects in the collision integral play an essential role in the description of the temperature dependence of the intrinsic width Γ. They lead to the saturated behavior of $\Gamma(T)$ as a function of T in contrast to the traditional Fermi-liquid theory, where the retardation effects are usually neglected; see Fig. 3 for the isovector GDR. The theoretical value of $\Gamma(T)$ at high temperatures is consistently smaller than the experimental data. We pointed out that a good agreement with experimental data can be obtained by adjusting the value of the parameter q in Eq. (21).

Our phenomenological model gives a reasonable explanation of the observed hindrance of the width of the isovector double giant dipole resonance (DGDR), see Fig. 5. The width of the DGDR is sensitive to the relative contributions of both the two-body and one-body dissipation mechanisms. Considering only the two-body dissipation, it is seen from Eqs. (21) and (24) that in the long relaxation time limit, i.e., zero-sound limit in cold nuclei, the memory effects in the collision integral lead to $\Gamma \sim E^2$. Therefore the width of the DGDR is larger by a factor of about 4 than the width of the single GDR because of the double eigenenergy $E = 2\hbar\omega_{1^-}$ of the DGDR. However, considering only the one-body dissipation channel we find that $\delta D = \Gamma_{\mathrm{double}}/\Gamma \approx 1$. The competition between these dissipation channels leads to the observable value and the A behavior of the ratio $\delta D = \Gamma_{\mathrm{double}}/\Gamma$ represented in Fig. 5.

The presence of both the one-body and particle emission contributions to the relaxation of the collective motion in a finite nucleus modifies the infinite nuclear matter condition for the transition from the zero-sound regime to the first-sound regime. We have noted that the bell-shaped form of the intrinsic width Γ as a function of $\omega\tau$ provides a new criterion in Eq. (61) for a determination of the transition temperature T_{tr} between the zero-sound and first-sound regimes in hot nuclei. This criterion is different from the case of infinite nuclear matter. In the case of the isovector GDR the corresponding transition temperature T_{tr} is significantly lower than the corresponding transition temperature in the infinite nuclear matter.

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