Isovector vibrations in nuclear matter at finite temperature

M. Di Toro,[1] V. M. Kolomietz,[2,3] and A. B. Larionov[1,4]
[1]*Dipartimento di Fisica and Istituto Nazionale di Fisica Nucleare, Laboratorio Nazionale del Sud, I-95123 Catania, Italy*
[2]*Institute for Nuclear Research, Prosp. Nauki 47, 252028 Kiev, Ukraine*
[3]*Cyclotron Institute, Texas A&M University, College Station, Texas 77843-3366*
[4]*Kurchatov Institute, Moscow 123182, Russia*
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We consider the propagation and damping of isovector excitations in heated nuclear matter within the Landau Fermi-liquid theory. Results obtained for nuclear matter are applied to calculate the giant dipole resonance (GDR) at finite temperature in heavy spherical nuclei within a Steinwedel and Jensen picture. The centroid energy of the GDR slightly decreases with increasing temperature and the width increases as T^2 for temperatures $T<4$ MeV in agreement with recent experimental data for GDR in ^{208}Pb and ^{120}Sn. The validity of the method for other Fermi fluids is finally suggested. [S0556-2813(99)00206-X]

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I. INTRODUCTION

In recent years the GDR built on highly excited states is a central issue of many experimental and theoretical studies (see Ref. [1], and references therein). In this context, one of the most important open problems is the behavior of the GDR width in *nonrotating* nuclei as a function of temperature. There are two essentially different theoretical approaches to this problem. The first one [2] explains the temperature increasing of the width as an effect of the adiabatic coupling of the GDR to thermal shape deformations. In the second approach [3–7] the thermal contribution to the damping width arises from an increasing nucleon-nucleon (NN) collision rate ($2p2h$ excitations) plus a Landau spreading due to new thermally allowed ph transitions [8–12].

In the present work, following the ideology of the second approach, we consider isovector volume vibrations in spin-isospin symmetrical nuclear matter at finite temperature. A similar problem was studied in Refs. [9,10] within the random phase approximation (RPA), one-body, method. However the Landau damping mechanism of the dissipation of a propagating mode due to the thermal smearing of a Fermi distribution is too weak to be responsible for the fast increase of the observed GDR width with temperature [9,10,13].

The puzzle can be solved by taking into account the two-body dissipation through the collision integral of the Landau-Vlasov equation [3]. We will use a quantum kinetic equation which leads to the introduction of memory effects in the collision term in order to include off energy-shell contributions [14]. Moreover, it was shown in Refs. [15,16], that memory effects are essentially increasing the widths of multipole resonances at small temperatures. In this work, we calculate the isovector strength function of nuclear matter taking into account both thermal Landau damping and two-body collisional dissipation, including the quantum memory contribution.

The structure of the work is as follows. In Sec. II the response function of the nuclear matter to the isovector external field is derived. Various dissipation mechanisms contributing to the damping width of the GDR in hot nuclei are discussed. In Sec. III the results of numerical calculations of the isovector strength functions are reported. The full width at half maximum (FWHM) of the theoretical photoabsorption cross section by heated ^{208}Pb and ^{120}Sn nuclei is compared to the experimental data. That gives a possibility to get some restrictions on the value of the in-medium NN scattering cross sections. Discussion and summary are given in Sec. IV.

II. RESPONSE FUNCTION

The isovector response of uniform nuclear matter is described by the linearized Landau-Vlasov equation with a collision term treated in the relaxation time approximation [4,11,16]

$$\frac{\partial}{\partial t}\delta f+\mathbf{v}\cdot\nabla_r\delta f-\nabla_r(\delta U+2\,\delta V)\cdot\nabla_p f_{\rm eq}=-\frac{1}{\tau}\,\delta f|_{l\geq 1}, \tag{1}$$

where $\mathbf{v}=\mathbf{p}/m^*$ is a velocity, $\delta f\equiv\delta f_n-\delta f_p$, $\delta U\equiv\delta U_n-\delta U_p$, and $\delta V\equiv(\delta V_n-\delta V_p)/2$ ($\delta V_q=\tau_q\delta V$, $\tau_n=+1$, $\tau_p=-1$) [9] are differences between neutron and proton distribution functions (DFs), mean fields, and external fields, respectively, $f_{eq}(\epsilon_p=p^2/2m^*)$ is the equilibrium finite temperature Fermi distribution, and the notation $l\geq 1$ means that the perturbation of the DF $\delta f|_{l\geq 1}$ in the collision integral includes only Fermi surface distortions with a multipolarity $l\geq 1$ in order to conserve the particle number in the collision processes [14]. The inclusion of the $l=1$ harmonic in the collision integral of Eq. (1), at variance with the isoscalar case [11], is due to nonconservation of the isovector current, i.e., due to a collisional friction force between counter-streaming neutron and proton flows.

Equation (1) is derived assuming isospin symmetric unperturbed nuclear matter, where isovector and isoscalar perturbations propagate independently. In the case of asymmetric nuclear matter, equations for isovector and isoscalar DF become coupled through terms $\propto(N-Z)/A$ (see Ref. [17]).

The dynamical component of the isovector mean field δU can be expressed in terms of the isovector Landau parameter F_0' :

$$\delta U=\frac{F_0'}{N(T)}\,\delta\rho, \tag{2}$$

where

$$\delta\rho(\mathbf{r};t)=\int\frac{g\,d\mathbf{p}}{(2\pi)^3}\,\delta f(\mathbf{r},\mathbf{p};t) \tag{3}$$

is the density perturbation, $g=2$ is the spin degeneracy factor, and

$$N(T)=\int\frac{g\,d\mathbf{p}}{(2\pi)^3}\left(-\frac{\partial f_{\mathrm{eq}}(\epsilon_p)}{\partial\epsilon_p}\right) \tag{4}$$

is the thermally averaged density of states, $N(0)=gp_Fm/2\pi^2$, where for simplicity we put $m^*=m=938$ MeV. It is assumed $\hbar=1$ everywhere in this work.

For an external field $\delta V\propto\exp(i\mathbf{kr}-i\omega t)$, periodic in space and time, the isovector collective response function [9] can be derived from Eq. (1):

$$\chi^{\mathrm{coll}}(\omega,\mathbf{k})=-\frac{\delta\rho}{\delta V}=\frac{2N(T)\chi_T^\tau}{1+F_0'\chi_T^\tau}, \tag{5}$$

where χ_T^τ is the intrinsic response function [18–20]. The explicit form of the function $\chi_T^\tau(\omega,\mathbf{k})$ is (details of derivation can be found in Ref. [11])

$$\chi_T^\tau(s)=-\frac{N(0)}{mp_FN(T)}\int_0^\infty dp\,\frac{p^2s\chi(\overline{p}s/p)}{s'+is''\chi(\overline{p}s/p)}\,\frac{\partial f_{\mathrm{eq}}(\epsilon_p)}{\partial\epsilon_p}, \tag{6}$$

where

$$\overline{p}=p_F\left(\frac{\overline{\epsilon}}{\epsilon_F}\right)^{1/2}, \tag{7}$$

$$\overline{\epsilon}=\frac{5}{3\rho_{\mathrm{eq}}}\int\frac{g\,d\mathbf{p}}{(2\pi)^3}\,\epsilon_pf_{\mathrm{eq}}(\epsilon_p), \tag{8}$$

$$\rho_{\mathrm{eq}}=\int\frac{g\,d\mathbf{p}}{(2\pi)^3}\,f_{\mathrm{eq}}(\epsilon_p) \tag{9}$$

are quasiparticle average momentum, average kinetic energy (normalized at $T=0$ on p_F and ϵ_F), and density, with the complex variable

$$s=s'+is'',\quad s''=\frac{m}{\tau\overline{p}k},\quad s'=\frac{\omega m}{\overline{p}k}, \tag{10}$$

$\chi(z)$ is a Legendre function of the second kind

$$\chi(z)=\frac{1}{2}\int_{-1}^{1}d\mu\,\frac{\mu}{\mu-z}. \tag{11}$$

Equation (6) for the intrinsic isovector response function has only a minor difference with the isoscalar case. Namely, to recover the isoscalar response function given by Eq. (30) in Ref. [11], one should change $is''\rightarrow is''(1+3s's\overline{p^2}/p^2)$ in the denominator of Eq. (6). This difference is just due to inclusion of the damping of the $l=1$ harmonic in the isovector channel. We note a misprint in the right-hand side (RHS) of Eq. (30) in Ref. [11]: it should be multiplied by a minus sign.

For a given momentum transfer k, the strength function per unit volume is

$$S_k(\omega)=\frac{1}{\pi}\mathrm{Im}(\chi^{\mathrm{coll}})=\frac{2N(T)\,\mathrm{Im}(\chi_T^\tau)/\pi}{[1+F_0'\,\mathrm{Re}(\chi_T^\tau)]^2+[F_0'\,\mathrm{Im}(\chi_T^\tau)]^2}. \tag{12}$$

The strength function satisfies the following energy weighted sum rule (EWSR) [9,10]:

$$\int_0^\infty d\omega\,\omega S_k(\omega)=\frac{k^2}{2m}\rho_0, \tag{13}$$

where $\rho_0=0.16\ \mathrm{fm}^{-3}$ is the nuclear saturation density.

Collective modes are given by the poles of the response function (5)

$$1+F_0'\chi_T^\tau(s)=0. \tag{14}$$

By solving Eq. (14) we obtain the complex frequency

$$\omega=\omega_R+i\omega_I=k\frac{\overline{p}}{m}(s-is''). \tag{15}$$

The application of the formalism discussed above to finite nuclei is based on the Steinwedel-Jensen (SJ) model [9,10,21] which describes the GDR in heavy nuclei as a volume polarization mode conserving the total density $\rho=\rho_n+\rho_p$. According to this model, we choose the wave number of the normal mode as $k=\pi/2R$, where R is the radius of a nucleus. Inside the nucleus, the unperturbed distribution of nucleons is supposed to be uniform. The SJ model gives a good overall reproduction of the ground state GDR energies for heavy spherical nuclei [22].

Equation (1) contains the two free parameters: the isovector Landau parameter F_0' and the relaxation time τ. The isovector Landau parameter F_0' at zero temperature can be expressed as a function of the symmetry energy coefficient β in the Weizsäcker mass formula as follows [23]:

$$F_0'(T=0)=\frac{3\beta}{\epsilon_F}-1. \tag{16}$$

For a Skyrme interaction, the coupling constant $f_0'\equiv F_0'(T)/N(T)$ has the following well-known structure (see Ref. [24]):

$$f_0'=\frac{p_F^2}{4}[t_2(2x_2+1)-t_1(2x_1+1)]-\frac{t_0}{2}(2x_0+1)-\frac{t_3}{12}(2x_3+1)\rho^\alpha. \tag{17}$$

We have chosen the T6 interaction of Ref. [25] for numerical calculations: $t_0=-1794.20$ MeV fm^3, $t_1=294.00$ MeV fm^5, $t_2=-294.00$ MeV fm^5, $t_3=12817.00$

MeV fm$^{3(1+\alpha)}$, $x_0=0.392$, $x_1=-0.500$, $x_2=-0.500$, $x_3=0.500$, $\alpha=1/3$. This interaction is momentum independent and gives exactly Eq. (2) for the mean field perturbation. Respectively, the term $\propto p_F^2$ in the RHS of Eq. (17) is equal to zero for the T6 interaction.

We see that even at the absence of the momentum dependence, the temperature still influences on the coupling constant through the the density ρ. According to Ref. [26], the mean square of a nuclear radius increases quadratically with temperature due to thermal pressure of nucleons

$$\langle r^2\rangle=\langle r^2\rangle_{T=0}(1+a_{r^2}T^2), \tag{18}$$

where $a_{r^2}=0.4\times10^{-2}$ MeV^{-2} for heavy nuclei. Therefore

$$\rho(T)\simeq\rho_0(1-a_\rho T^2),\quad a_\rho=\frac{3}{2}a_{r^2}. \tag{19}$$

For the coupling constant f_0' and the Landau parameter F_0', we obtain the following expressions:

$$f_0'(T)=f_0'(0)(1+a_f T^2), \tag{20}$$

$$F_0'(T)=F_0'(0)(1+a_F T^2), \tag{21}$$

where $f_0'(0)$ is defined by Eq. (17) with $\rho=\rho_0$, and

$$a_f=\frac{t_3}{12f_0'(0)}(2x_3+1)\rho_0^\alpha\alpha a_\rho,$$

$$a_F=a_f-\frac{\pi^2}{12\epsilon_F^2}.$$

In the derivation of Eq. (21) we also used the low-temperature expansion of the level density

$$N(T)=N(0)\left(1-\frac{\pi^2}{12}\frac{T^2}{\epsilon_F^2}\right). \tag{22}$$

The thermal dependence of $N(T)$ is very weak with respect to the one of $f_0'(T)$ and can be neglected. That gives for the interaction T6 $a_F\simeq a_f\simeq0.5\times10^{-2}$ MeV^{-2}. The zero-temperature isovector Landau parameter has a value $F_0'(0)=1.43$ corresponding to the symmetry energy $\beta=30$ MeV and $\epsilon_F=37$ MeV [see Eq. (16)].

The relaxation time τ generally includes various dissipation mechanisms (see Ref. [16]):

$$\tau^{-1}=\tau_{\rm coll}^{-1}+\tau_{\rm wall}^{-1}+\tau_\uparrow^{-1}, \tag{23}$$

where $\tau_{\rm coll}$, $\tau_{\rm wall}$, and $\tau_\uparrow$ are the relaxation times due to two-body collisions, wall friction, and particle emission, respectively. The most important dissipation mechanism, which defines the temperature trend of the total width, is the dissipation due to NN collisions. The corresponding two-body relaxation time $\tau_{\rm coll}$ includes the temperature and memory effects:

$$\tau_{\rm coll}=\frac{\alpha^{(-)}}{T^2+(\omega_R/2\pi)^2}. \tag{24}$$

The dependence of the relaxation time on the frequency ω_R arises from the memory effects and corresponds to the Landau prescription [14]. The coefficient $\alpha^{(-)}$ depends on the NN scattering cross sections. We have calculated this coefficient using energy and angular dependent differential cross sections of pp and np scattering derived from the Bonn A potential both with and without in-medium corrections [27,28] (see the Appendix). The results are $\alpha^{(-)}=2.3$ (5.4) MeV in the case of vacuum (in-medium reduced) cross sections.

The wall friction is related to the fragmentation width [29], i.e., the GDR spreading due to a coupling to ph excitations. It can be taken into account as an additional dissipative source term in the kinetic equation with a relaxation time (see Ref. [16])

$$\tau_{\rm wall}=\frac{2R}{\bar{v}}\xi, \tag{25}$$

where R is a nuclear radius,

$$\bar{v}=\frac{3v_F}{4}\left[1+\frac{\pi^2}{6}\left(\frac{T}{\epsilon_F}\right)^2\right]$$

is the temperature-dependent average velocity of nucleons, and $\xi\sim10$ is a numerical factor. Actually the value of this factor can be chosen to fit the fragmentation width of the GDR in a cold nucleus, which is of the order of 0.5–2 MeV as given by RPA studies, Refs. [30,31]. Since the temperature region of interest is $T\ll\epsilon_F$, the relaxation time $\tau_{\rm wall}$ depends on temperature very weakly. This fact is indeed in agreement with a weak temperature dependence of the fragmentation width obtained in RPA calculations, Ref. [32].

The relaxation time $\tau_\uparrow$ caused by the direct particle emission is related to the escape width $\Gamma_\uparrow$ (see Ref. [33]). From semiclassical [3,33] as well as quantal continuum RPA calculations [34] this quantity turns out to be quite small at all temperatures. The reason for this is the small amplitude of the momentum distortions which leads to a quite reduced probability of direct particle emission. Therefore we will neglect the contribution of particle escape to the total relaxation time τ.

III. NUMERICAL RESULTS

The analysis of a hot GDR is based on the photoabsorption cross section by a thermally excited nucleus $\sigma_{\rm abs}(\omega)$ which can be expressed in terms of the strength function $S_k(\omega)$ of Eq. (12) as follows:

$$\sigma_{\rm abs}(\omega)=\frac{4\pi^2e^2}{ck^2\rho_0}\frac{NZ}{A}\omega S_k(\omega). \tag{26}$$

This expression is obtained from comparison of the EWSR (13) and the dipole sum rule of Ref. [21]:

$$\int_0^\infty dE\,\sigma_{\rm abs}(E)=\frac{2\pi^2e^2}{mc}\frac{NZ}{A}. \tag{27}$$

In this work we have studied the photoabsorption cross section at temperatures $T\leqslant4$ MeV using vacuum and in-

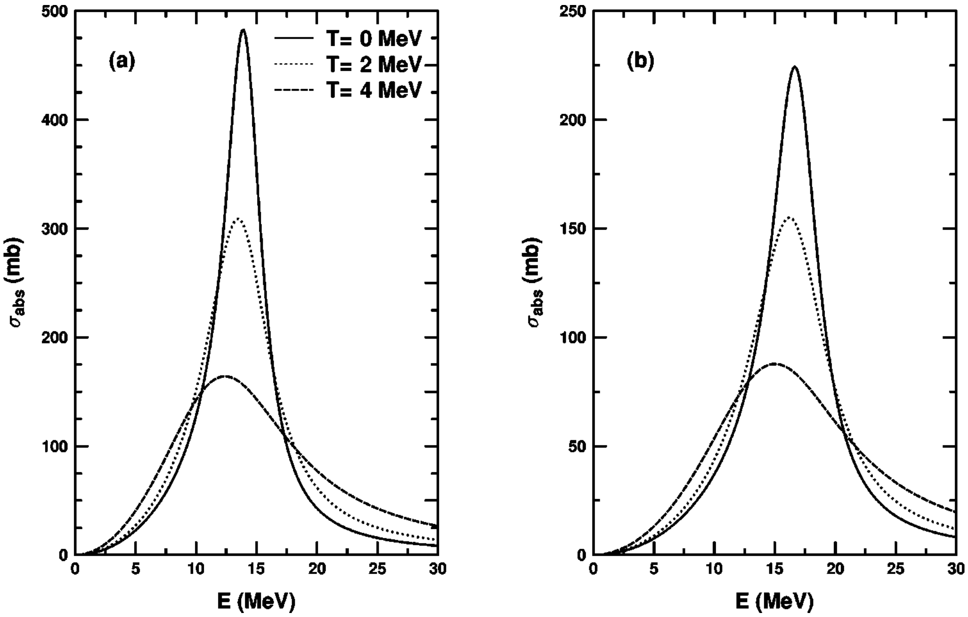

FIG. 1. Photoabsorption cross section by an excited nucleus ^{208}Pb (a) and by ^{120}Sn (b) at various temperatures as a function of the photon energy. Calculations are done using vacuum cross sections plus a wall friction contribution fitted to the g.s. FWHM.

medium reduced NN scattering cross sections. Calculations with a somewhat more developed version of our collisional damping approach, including different relaxation times for harmonics with $l=1$ and $l \geqslant 2$, have been done, actually, for temperatures up to 10 MeV to eventually see a zero-to-first sound transition. These results will be presented in a forthcoming paper, Ref. [35].

In both cases of vacuum and in-medium reduced cross sections, we have performed calculations switching on and off the wall friction contribution $\tau_{\rm wall}^{-1}$ in the inverse relaxation time τ^{-1} [see Eq. (23)]. A value of this contribution was chosen to fit the GDR width in ground state (g.s.) nuclei. Of course when we use in-medium reduced cross sections we need a larger one-body damping. The following values of the wall-friction width $\Gamma_{\rm wall}=2/\tau_{\rm wall}$ have been obtained: $\Gamma_{\rm wall}=0.5$ (3.2) MeV for ^{208}Pb and $\Gamma_{\rm wall}=0.9$ (3.8) MeV for ^{120}Sn for vacuum (in-medium reduced) NN scattering cross sections. Once the ratio wall/collisional damping at $T=0$ is fixed, one can predict unambiguolsy the temperature behavior of the total width. Below, if the opposite is not specially indicated, the discussion is done for the case of vacuum NN cross sections plus wall-friction contribution.

Figure 1 shows the photoabsorption cross section by nuclei ^{208}Pb (a) and ^{120}Sn (b) as a function of the photon energy for several temperatures. As the temperature is growing, the centroid energy $E_{\rm GDR}$ (i.e., the peak of the photoabsorption cross section) is shifting to smaller energies and the width is increasing. In Fig. 2 (solid line) we report the temperature dependence of the corresponding centroid energy for ^{208}Pb. To understand better the thermal behavior of $E_{\rm GDR}$, we present in Fig. 2 (see dotted and dashed lines) also the calculations performed for a fixed nuclear radius R (T independent) in the formula for the wave number $k=\pi/2R$ of the normal mode. We see that $E_{\rm GDR}$ decreases mostly due to the thermal expansion of a nucleus. However, it slightly decreases even for a fixed radius R (see dotted line in Fig. 2). Indeed at larger temperatures an increasing two-body dissipation should reduce the frequency of the collective motion, in close analogy with a classical oscillator with a friction force (see also Refs. [36–38]). This is at variance with pure mean field predictions [9–11] of a growth of $E_{\rm GDR}$ with temperature. To conclude the discussion on the centroid energy, we note, that the thermal dependence of the coupling constant f_0', introduced according to Eq. (20) (see dashed line in Fig. 2) does not, practically, change the results obtained at constant f_0' (dotted line).

In Fig. 3 (solid lines), the real and imaginary parts of the pole of the response function (5) are shown as functions of temperature for a nucleus ^{208}Pb. As far as the collective mode is underdamped, i.e., $|{\rm Im}(\omega)|/{\rm Re}(\omega) \ll 1$, an approximate formula

$$\mathrm{FWHM}=2|\mathrm{Im}(\omega)| \tag{28}$$

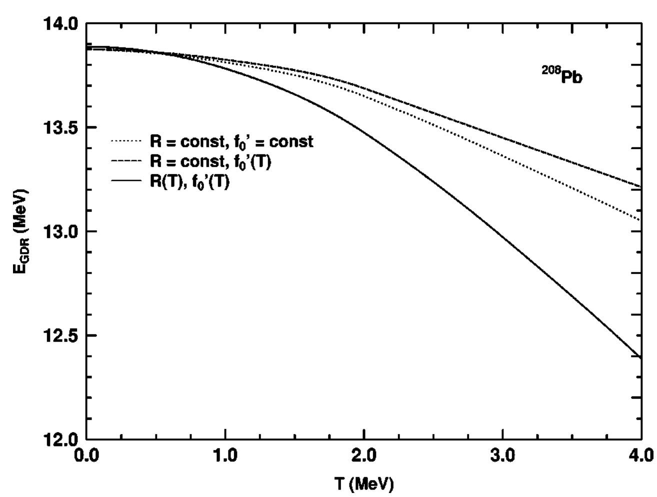

FIG. 2. Temperature dependence of the GDR energy $E_{\rm GDR}$ defined at the peak position of the photoabsorption cross section for a nucleus ^{208}Pb [see Fig. 1(a)]. Dotted, dashed, and solid lines correspond to calculations with fixed radius R and fixed isovector coupling constant f_0', with fixed R and temperature dependent f_0', and with temperature dependent R and f_0' [see Eqs. (18), (20)]. A weak temperature dependence of the level density $N(T)$ [see Eq. (22)] is always present in the Landau parameter $F_0'=N(T)f_0'$. Vacuum cross sections plus wall friction contribution are used.

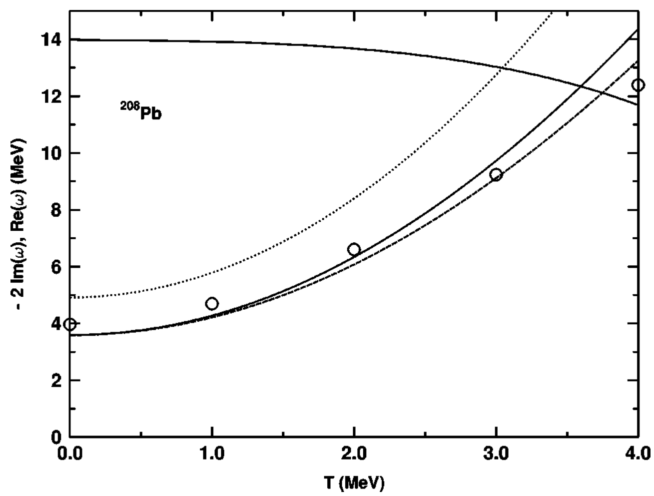

FIG. 3. Temperature dependence of $\mathrm{Re}(\omega)$, where ω is the pole of the response function (5) (upper solid line) and of $-2\mathrm{Im}(\omega)$, where $\mathrm{Im}(\omega)$ is the imaginary part of the pole of the response function (5) (lower solid line) defined by Eq. (32) (dashed line) defined by Eq. (34) (dotted line.) Open circles show the FWHM extracted from the photoabsorption cross section of Fig. 1(a). Calculations are for a nucleus ^{208}Pb using vacuum cross sections plus wall friction contribution.

has to be fulfilled [16]. We see from Fig. 3 (compare lower solid line to the open circles), that the condition (28) is really satisfied, corresponding to the Lorentzian shape of the photoabsorption strength at high temperatures (see Fig. 1 and Ref. [39]). Some relatively small deviations from Eq. (28) present in Fig. 3 can be explained by a quite large width already at $T=0$. For instance, we have checked that Eq. (28) is fulfilled with a much better accuracy in the case of in-medium reduced NN cross sections without wall friction contribution.

At low temperatures, one can neglect the temperature spreading of the equilibrium Fermi distribution substituting $\delta(\epsilon_F-\epsilon_p)$ instead of $[-\partial f_{\mathrm{eq}}(\epsilon_p)/\partial\epsilon_p]$ into Eq. (6). Thus we obtain the following approximate low-temperature expression for the intrinsic response function:

$$\chi_T^\tau(s)\simeq\frac{s\chi(s)}{s'+is''\chi(s)}, \tag{29}$$

where variables s, s'', and s' are defined by Eqs. (10) with the change $\overline{p}\to p_F$. We remark that Eq. (29) when applied to the case of an electron gas $[F_0=N(T)4\pi e^2/k^2]$ gives just the longitudinal dielectric function

$$\epsilon(\mathbf{k},\omega)=1+F_0\chi_T^\tau \tag{30}$$

obtained by Mermin [40]. This is to stress the more general framework of the results presented here [41].

In the rare collision regime $(\omega_R\tau\gg 1)$, an approximate solution of the dispersion relation (14) with the intrinsic response function (29) can be found analytically (see Refs. [14,42]):

$$\omega_R\simeq v_F k s^{(0)}+O(T^4), \tag{31}$$

$$\omega_I\simeq\frac{(2F_0'+1)[(s^{(0)})^2-1]-(F_0')^2}{\tau F_0'[F_0'-(s^{(0)})^2+1]}+O(T^6), \tag{32}$$

where $s^{(0)}$ is the root of collisionless dispersion relation

$$1+F_0'\chi(s^{(0)})=0. \tag{33}$$

The Landau parameter F_0' in Eqs. (32), (33) is taken at $T=0$. The simple expression (32) for the imaginary part of the frequency ω (dashed line in Fig. 3) reproduces the results of a numerical solution of the ''exact'' dispersion relation (14) (solid line in the same Fig. 3) with a good accuracy for temperatures $T<2$ MeV. At larger temperatures, a slight increase of the damping due to temperature smearing of the Fermi distribution is obtained with the dispersion relation (14). The difference between these two solutions is of the order of the Landau damping rate in the pure mean field approach [9,10]. We see that thermal Landau damping is small for the case of GDR, at variance with the case of the isoscalar mode [11], since the isovector Landau parameter is larger than the isoscalar one for nuclear effective interactions at normal density. As a consequence a weaker coupling between single particle and collective motion is expected for isovector vibrations [19].

The dominance of the collisional contribution to the total damping rate of the GDR is just expressed by an approximate relation (dotted line in Fig. 3)

$$-\mathrm{Im}(\omega)\simeq\frac{1}{\tau}, \tag{34}$$

which was used, for instance, in Ref. [15] to calculate widths of giant resonances. The main deviation from Eq. (34) in the dispersion relation (14) is caused by the exclusion of the $l=0$ harmonic from the collision integral on the RHS of Eq. (1), i.e., due to taking into account the particle number conservation. This results in a smaller absolute value of $\mathrm{Im}(\omega)$.

A source of uncertainty in our calculations is given by the choice of the NN cross sections [see Eq. (A23) in the Appendix]. This is caused mostly by two reasons (i) The density in a finite nucleus is not uniform. It is clear that vacuum cross sections are more appropriate in the surface region. (ii) A quite large uncertainty of the in-medium NN cross sections themselves exists, as a result from various calculations (see Refs. [27,43]).

An explicit implementation of the density-dependent in-medium cross section for the space nonuniform density distribution of a finite nucleus, as was done in Ref. [44], is out of the scope of the present work. Thus, we have considered the two limiting cases of vacuum cross sections and in-medium reduced cross sections of Refs. [27,28] at the nuclear saturation density ρ_0. In Figs. 4(a)–4(d) calculations are presented with vacuum NN cross sections (a), (b) and with in-medium ones (c), (d) in comparison to the experimental widths from Refs. [45,46] for the nuclei ^{208}Pb (a), (c) and ^{120}Sn (b), (d). The use of vacuum cross sections gives much better agreement with experiment compared to the case of in-medium reduced cross sections. This conclusion is not influenced by the switching on/off the wall friction contribution [compare solid and dotted lines in Figs. 4(a)–4(d)]. We see from Figs. 4(c), 4(d), that the wall friction contribu-

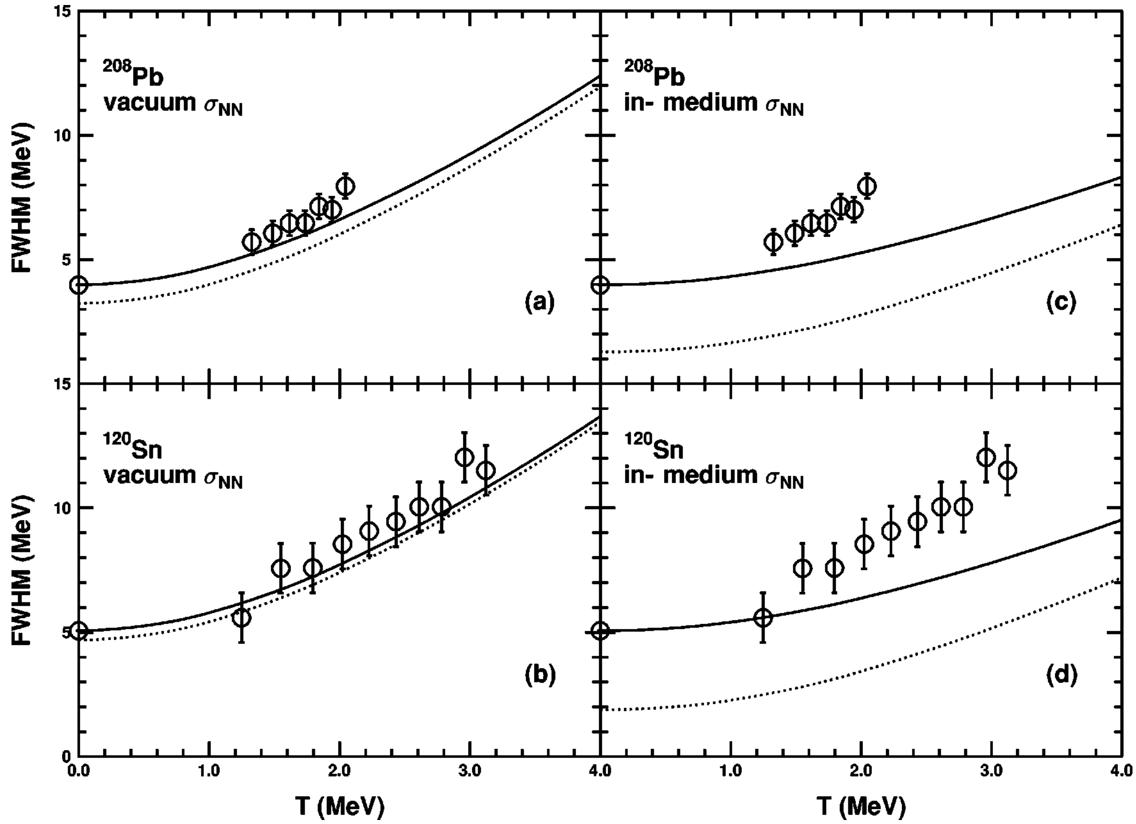

FIG. 4. The width FWHM of the GDR mode in nuclei ^{208}Pb (a), (c) and ^{120}Sn (b), (d) as a function of temperature. Solid and dotted lines represent calculations with and without a wall friction contribution, respectively. Panels (a), (b) [(c), (d)] contain results obtained with vacuum (in-medium reduced) NN scattering cross sections. Points with errorbars show the experimental widths from Ref. [45] (^{208}Pb) and from Ref. [46] (^{120}Sn).

tion is important mostly at low temperature and gradually decreases at high temperature. That is, in spite of the increasing one-body damping rate $\tau_{\mathrm{wall}}^{-1}$ with temperature [see Eq. (25)] and due to the nonlinear deviations from Eq. (34), caused by imposing the particle number conservation on the RHS of Eq. (1).

IV. DISCUSSION AND SUMMARY

We have studied the isovector response of a heated spin-isospin symmetric nuclear matter on the basis of the linearized Landau-Vlasov equation with a collision integral including memory effects in the relaxation time approximation. Further, the dipole polarization SJ mode in a finite nucleus has been considered.

The one-body dissipation due to wall friction has been included as an additional source term in the kinetic equation that simulates the fragmentation width contribution (see Refs. [16,29]). Since the aim of this work was to study the thermal behavior of the GDR damping width, we have fitted a zero-temperature experimental width varying a value of the wall friction contribution as a free parameter and keeping the collisional contribution fixed. Once this fit is done, the temperature dependence of the width is defined unambiguously for a given set of the NN scattering cross sections.

We have found that vacuum NN cross sections, used neglecting effective mass effects, give better agreement with experimental data on the GDR width at finite temperature. That implies a quite small (~ 0.5 MeV) contribution of the fragmentation width to the total GDR damping width already at $T=0$ [see Figs. 4(a), 4(b)]. The thermal Landau damping due to the coupling of thermally excited quasiparticles to the collective mode is also found to be negligible. Thus the relaxation of the volume isovector mode is caused mainly by NN collisions. An increasing temperature shifts the centroid energy of the isovector strength function to smaller values, mostly due to the thermal swelling of a nucleus. The calculated width is proportional to $1/\tau$ in the temperature region $T<4$ MeV studied here, that corresponds to a dominant collisional damping of the isovector zero sound mode and leads to a T^2 behavior of the GDR-FWHM.

We have shown that some general GDR properties at high excitation energy can be obtained directly from Fermi-liquid theory. However, the main purpose of this work is not to get perfect agreement with data for finite nuclei. Indeed, we are well aware that other contributions to the damping are missing in the present approach: (i) Thermal shape fluctuations [2,47]; (ii) fluctuations due to NN correlations [48]. In particular the inclusion of these contributions could improve the agreement with the data in the case of in-medium reduced NN cross sections [see Figs. 4(c), 4(d)]. Moreover, the inclusion of the effective mass corrections in the calculation of the collisional relaxation times Eq. (A12) is expected to reduce the collisional contribution to the total width (see Ref. [44]). Therefore, some room for the contributions (i),(ii) will also appear in the case of vacuum cross sections [see Figs. 4(a), 4(b)].

A recent statistical model analysis of γ spectra produced by inelastic α scattering on ^{120}Sn [39] resulted in the conclusion that neither the thermal fluctuation model of Ref. [47] nor the collisional damping model could reproduce data in details. We believe that the combination of the two models would give a much better agreement in that low-temperature region, where the adiabatic coupling to shape fluctuations is more justified.

Extension to isospin asymmetric nuclear systems, more suitable for the Pb case, can be performed following the approach of Refs. [17,49,50]. We do not expect to have substantial variations in the temperature behavior of the isovector and isoscalar modes unless very high charge asymmetries are reached [17,49]. However, in charge asymmetric nuclei, a new soft mode different than the isovector and isoscalar ones seems to appear [50] due to the collisional coupling of proton and neutron vibrations, that requires further investigation.

Finally, a comment on the validity of the formalism applied here to the study of hot GDR in nuclei to a broader context of physical problems. Indeed it can be easily generalized to the description of relative vibrations of any two-component Fermi liquid with a mutual attraction, for instance, of a Coulomb plasma consisting of opposite charged fermions. Another application could be to the oscillations of the electronic cloud in metallic clusters, where the momentum nonconserving $l=1$ term in the collision integral appears due to scattering of electrons on impurities [40].

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APPENDIX

Here we will calculate the coefficient $\alpha^{(-)}$ in Eq. (24) for the collisional relaxation time $\tau_{\rm coll}$, closely following the formalism of Refs. [51,52]. For simplicity, we will use in the derivation the Boltzmann-Uehling-Uhlenbeck (BUU) collision integrals without memory effects. However, as is shown in Ref. [44], the collision integrals with memory effects give the result which can be obtained using Landau prescription $\tau_{\rm coll}=\tau_{\rm BUU}/[1+(\omega_R/2\pi T)^2]$, where $\tau_{\rm BUU}=\alpha^{(-)}/T^2$ is the relaxation time given by the BUU collision integrals.

Time evolution of the space-uniform isovector DF $f=f_n-f_p$ satisfies the equation (see Refs. [15,16])

$$\frac{\partial f(\mathbf{p};t)}{\partial t}=I=I_{nn}+I_{np}-I_{pp}-I_{pn}, \tag{A1}$$

where $I_{q_1q_2}$ stands for the collision integral of particles of the sort $q_1=n,p$ with particles of the sort $q_2=n,p$. Explicitly

$$\begin{aligned}I_{q_1q_2}(\mathbf{p_1};t)=&\frac{g^2}{(2\pi)^6}\int d\mathbf{p_2}d\mathbf{p_3}d\mathbf{p_4}\\&\times w_{q_1q_2}(\mathbf{p_1},\mathbf{p_2};\mathbf{p_3},\mathbf{p_4})\,\delta(\triangle\epsilon)\,\delta(\triangle\mathbf{p})Q\\&\times[f_{q_1}(\mathbf{p_1};t),f_{q_2}(\mathbf{p_2};t);f_{q_1}(\mathbf{p_3};t),f_{q_2}(\mathbf{p_4};t)],\end{aligned} \tag{A2}$$

where

$$Q(f_1,f_2;f_3,f_4)\equiv(1-f_1)(1-f_2)f_3f_4-f_1f_2(1-f_3)(1-f_4),$$

$w_{q_1q_2}(\mathbf{p_1},\mathbf{p_2};\mathbf{p_3},\mathbf{p_4})$ is the spin-averaged probability of two-body collisions with initial momenta $(\mathbf{p_1},\mathbf{p_2})$ and final momenta $(\mathbf{p_3},\mathbf{p_4})$, $\triangle\epsilon=\epsilon_{p_1}+\epsilon_{p_2}-\epsilon_{p_3}-\epsilon_{p_4}$, $\triangle\mathbf{p}=\mathbf{p_1}+\mathbf{p_2}-\mathbf{p_3}-\mathbf{p_4}$. Neglecting the dependence of $w_{q_1q_2}$ on the DF, we can write down perturbations of collision integrals (A2) keeping the terms of the first order in δf_q:

$$\begin{aligned}\delta I_{q_1q_2}(\mathbf{p_1};t)=&\frac{g^2}{(2\pi)^6}\int d\mathbf{p_2}d\mathbf{p_3}d\mathbf{p_4}\\&\times w_{q_1q_2}(\mathbf{p_1},\mathbf{p_2};\mathbf{p_3},\mathbf{p_4})\,\delta(\triangle\mathbf{p})\{\alpha^{(1)}\psi_{q_1}(\mathbf{p_1};t)\\&+\alpha^{(2)}\psi_{q_2}(\mathbf{p_2};t)+\alpha^{(3)}\psi_{q_1}(\mathbf{p_3};t)\\&+\alpha^{(4)}\psi_{q_2}(\mathbf{p_4};t)\},\end{aligned} \tag{A3}$$

where

$$\psi_q(\mathbf{p};t)\equiv\delta f_q(\mathbf{p};t)\left(\frac{\partial f_{eq}(\epsilon_p)}{\partial\epsilon_p}\right)^{-1}\quad(q=n,p), \tag{A4}$$

$$\begin{aligned}\alpha^{(i)}\equiv&\frac{\delta Q[f_{eq}(\epsilon_{p_1}),f_{eq}(\epsilon_{p_2});f_{eq}(\epsilon_{p_3}),f_{eq}(\epsilon_{p_4})]}{\delta f_{eq}(\epsilon_{p_i})}\\&\times\frac{\partial f_{eq}(\epsilon_{p_i})}{\partial\epsilon_{p_i}}\,\delta(\triangle\epsilon),\quad i=1,2,3,4.\end{aligned} \tag{A5}$$

In Eq. (A3) the isospin-symmetric nuclear matter is considered that results in the same equilibrium DF for neutrons and protons.

For the perturbation of the collision integral I of the RHS of Eq. (A1), assuming isotopic invariance ($w_{pp}=w_{nn}$, $w_{pn}=w_{np}$) we can write after simple algebra

$$\begin{aligned}\delta I(\mathbf{p_1};t)=&\delta I_{nn}+\delta I_{np}-\delta I_{pp}-\delta I_{pn}\\=&\frac{g^2}{(2\pi)^6}\int d\mathbf{p_2}d\mathbf{p_3}d\mathbf{p_4}\delta(\triangle\mathbf{p})\{(w_{pp}+w_{np})\\&\times[\alpha^{(1)}\psi(\mathbf{p_1};t)+\alpha^{(3)}\psi(\mathbf{p_3};t)]+(w_{pp}-w_{np})\\&\times[\alpha^{(2)}\psi(\mathbf{p_2};t)+\alpha^{(4)}\psi(\mathbf{p_4};t)]\},\end{aligned} \tag{A6}$$

where $\psi(\mathbf{p};t)=\psi_n(\mathbf{p};t)-\psi_p(\mathbf{p};t)$. The triple integral over momenta in Eq. (A6) can be taken using Abrikosov-Khalatnikov transformation (see Ref. [14]), which is valid in the limit $T\ll\epsilon_F$:

$$\int d\mathbf{p_2} d\mathbf{p_3} d\mathbf{p_4} \delta(\triangle \mathbf{p}) = \frac{(m^*)^3}{2} \int_0^\pi d\Theta \frac{\sin\Theta}{\cos(\Theta/2)} \times \int_0^\pi d\phi \int_0^{2\pi} d\phi_2 \int_0^\infty d\epsilon_2 d\epsilon_3 d\epsilon_4, \tag{A7}$$

where $\theta = \widehat{(\hat{p}_1, \hat{p}_2)}$ is the angle between momenta of colliding particles ($\hat{p}_i \equiv \mathbf{p_i}/|\mathbf{p_i}|$, $i = 1,2,3,4$), ϕ is the angle between the planes given by the momenta of incoming and outcoming particles

$$\cos\phi = \frac{[\hat{p}_1 \times \hat{p}_2] \cdot [\hat{p}_3 \times \hat{p}_4]}{|[\hat{p}_1 \times \hat{p}_2]||[\hat{p}_3 \times \hat{p}_4]|}, \tag{A8}$$

ϕ_2 is the azimutal angle of $\mathbf{p_2}$ in the system with the z axis along $\mathbf{p_1}$.

We decompose the perturbation of the DF into spherical harmonics

$$\psi(\mathbf{p}_i;t) = \sum_{l,m} \alpha_{lm}(p_i;t) Y_{lm}(\hat{p}_i), \quad i = 1,2,3,4, \tag{A9}$$

where coefficients α_{lm} can be taken on the Fermi surface, since $T \ll \epsilon_F$. According to Refs. [15,42,53], the partial relaxation time τ_l^{BUU} is defined as follows:

$$\frac{1}{\tau_l^{\mathrm{BUU}}} = -\frac{\int_0^\infty d\epsilon_p \int d\Omega_{\hat{p}} Y^*_{lm}(\hat{p}) \delta I(\mathbf{p};t)}{\int_0^\infty d\epsilon_p \int d\Omega_{\hat{p}} Y^*_{lm}(\hat{p}) \delta f(\mathbf{p};t)} = \frac{\int d\Omega_{\hat{p}} Y^*_{lm}(\hat{p}) \overline{\delta I(\hat{p};t)}}{\alpha_{lm}}, \tag{A10}$$

where

$$\overline{\delta I(\hat{p};t)} = \int_0^\infty d\epsilon_p \delta I(\mathbf{p};t). \tag{A11}$$

Using Eq. (A7), after somewhat lengthy but standard calculations, we come to the expression

$$\frac{1}{\tau_l^{\mathrm{BUU}}} = \frac{(m^*)^3 T^2}{12\pi^2} \{\langle w_{pp} \Phi_l^{(+)} \rangle + 2 \langle w_{np} \Phi_l^{(-)} \rangle\} \equiv \frac{T^2}{\kappa_l}, \tag{A12}$$

where angular brackets denote the averaging over angles θ and ϕ [16,42,51]

$$\langle F(\theta,\phi) \rangle \equiv \frac{1}{2\pi} \int_0^\pi d\theta \frac{\sin\theta}{\cos\frac{\theta}{2}} \int_0^\pi d\phi F(\theta,\phi), \tag{A13}$$

$$\Phi_l^{(+)} = 1 + P_l(\hat{p}_2 \cdot \hat{p}_1) - P_l(\hat{p}_3 \cdot \hat{p}_1) - P_l(\hat{p}_4 \cdot \hat{p}_1), \tag{A14}$$

$$\Phi_l^{(-)} = 1 - P_l(\hat{p}_2 \cdot \hat{p}_1) - P_l(\hat{p}_3 \cdot \hat{p}_1) + P_l(\hat{p}_4 \cdot \hat{p}_1). \tag{A15}$$

A factor of 2 in the second term in curly brackets of Eq. (A12) is due to half momentum space integration over $d\mathbf{p_3}$ on the LHS of Eq. (A7) (see Refs. [14,54,55]). The arguments of the Legendre polynomials in functions (A14), (A15) are

$$\hat{p}_2 \cdot \hat{p}_1 = \cos\theta,$$

$$\hat{p}_3 \cdot \hat{p}_1 = \cos^2\frac{\theta}{2} + \sin^2\frac{\theta}{2}\cos\phi,$$

$$\hat{p}_4 \cdot \hat{p}_1 = \cos^2\frac{\theta}{2} - \sin^2\frac{\theta}{2}\cos\phi.$$

For $l = 1,\ 2$, and ∞ we have

$$\Phi_1^{(+)} = 0, \quad \Phi_1^{(-)} = 4\sin^2\frac{\theta}{2}\sin^2\frac{\phi}{2}, \tag{A16}$$

$$\Phi_2^{(+)} = 3\sin^4\frac{\theta}{2}\sin^2\phi, \quad \Phi_2^{(-)} = 3\sin^2\theta\sin^2\frac{\phi}{2}, \tag{A17}$$

$$\Phi_\infty^{(+)} = \Phi_\infty^{(-)} = 1. \tag{A18}$$

Collision probabilities can be expressed in terms of cross sections as follows:

$$w_{pp} = \frac{(2\pi)^3}{2\mu^2} \frac{d\sigma_{pp}}{d\Omega_{\mathrm{c.m.}}}, \tag{A19}$$

$$w_{pn} = \frac{(2\pi)^3}{2\mu^2} \frac{d\sigma_{pn}}{d\Omega_{\mathrm{c.m.}}}, \tag{A20}$$

where $d\sigma_{pp}/d\Omega_{\mathrm{c.m.}}$ and $d\sigma_{pn}/d\Omega_{\mathrm{c.m.}}$ are differential cross sections of pp and np scattering, $d\Omega_{\mathrm{c.m.}} = \sin\theta_{\mathrm{c.m.}} d\theta_{\mathrm{c.m.}} d\phi_{\mathrm{c.m.}}$, $\theta_{\mathrm{c.m.}}$, and $\phi_{\mathrm{c.m.}}$ are polar and azimutal scattering angles in the center of mass system of colliding particles, $\mu = m^*/2$ is the reduced mass. Differential cross sections $d\sigma_{pp}/d\Omega_{\mathrm{c.m.}}$ and $d\sigma_{pn}/d\Omega_{\mathrm{c.m.}}$ depend on the relative momentum $p' = |\mathbf{p_1} - \mathbf{p_2}|/2$ of scattered particles and on the polar angle

TABLE I. Parameters κ_l, $l = 1,\ 2$, and ∞ (MeV^2 fm/c) defined in Eq. (A12) at various choices of nucleon-nucleon scattering cross sections.

NN cross sections	κ_1	κ_2	κ_∞
Vacuum of Ref. [28]	503	491	401
In-medium at $\rho = \rho_0$ of Ref. [28]	1123	1068	1003
Isotropic energy-independent of Ref. [15]	920	1022	818

$$\theta_{\text{c.m.}} = \arccos\left(\frac{(\mathbf{p_1}-\mathbf{p_2})\cdot(\mathbf{p_3}-\mathbf{p_4})}{|\mathbf{p_1}-\mathbf{p_2}||\mathbf{p_3}-\mathbf{p_4}|}\right).$$

For particles scattered on the Fermi surface, we have

$$p' = p_F \sin\frac{\theta}{2}, \quad \theta_{\text{c.m.}} = \phi.$$

In a particular case of isotropic energy-independent cross sections the result of Ref. [15] is recovered:

$$\frac{1}{\tau_1^{\text{BUU}}} = \frac{32}{9} m \sigma_v T^2, \tag{A21}$$

$$\frac{1}{\tau_2^{BUU}} = \frac{32}{15} m \sigma_s T^2, \tag{A22}$$

where $\sigma_v = \sigma_{np}/2$, $\sigma_s = (\sigma_{nn} + \sigma_{pp} + 2\sigma_{np})/4$, $\sigma_{np} = (4\pi) d\sigma_{np}/d\Omega_{\text{c.m.}} \simeq 50$ mb, $\sigma_{nn} \simeq \sigma_{pp} = (2\pi) d\sigma_{pp}/d\Omega_{c.m.} \simeq 25$ mb.

We derived relaxation times τ_1^{BUU}, τ_2^{BUU} and τ_∞^{BUU} using pp and nn energy- and angular-dependent differential cross sections calculated with Bonn A potential in Refs. [27,28] and putting $m^* = m$ in Eq. (A12). Results of these calculations both with vacuum and in-medium cross sections at normal nuclear density are given in Table I. It is seen from the table that always $\tau_1^{\text{BUU}} \simeq \tau_2^{\text{BUU}} \simeq \tau_\infty^{\text{BUU}}$. That gives the idea of putting the same value for all relaxation times τ_l^{BUU}, i.e., to apply the usual relaxation time approximation. Thus, we obtain the collision integral of Eq. (1) with the relaxation time τ_{coll} given by Eq. (24), where

$$\alpha^{(-)} = \frac{3}{\kappa_1^{-1} + \kappa_2^{-1} + \kappa_\infty^{-1}} = 2.3 \ (5.4) \ \text{MeV}, \tag{A23}$$

for vacuum (in-medium) cross sections of Ref. [28].

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