Vol. 39 (2008) *ACTA PHYSICA POLONICA B* No 6

ISOVECTOR RESPONSE OF NUCLEAR MATTER AT FINITE TEMPERATURE

S. Ayik

Physics Department, Tennessee Technological University
Cookeville, TN 38505, USA

K. Bozkurt

Physics Department, Yildiz Technical University
34010 Esenler, Istanbul, Turkey

A. Gokalp, O. Yilmaz

Physics Department, Middle East Technical University
06531 Ankara, Turkey

(Received January 28, 2008)

The dipole response function of nuclear matter at zero and finite temperatures is investigated in an extended RPA approach by including collisional damping mechanism and coherent damping due to particle–phonon coupling. Calculations are carried out for nuclear dipole vibrations by employing the Steinwedel–Jensen model and compared with experimental results for ^{120}Sn and ^{208}Pb.

PACS numbers: 21.60.Jz, 21.65.+f, 24.30.Cz, 25.70.Lm

Studies of giant resonance excitations in atomic nuclei, in particular isovector giant dipole resonance (GDR), have been the subject of many experimental and theoretical studies [1]. A large amount of experimental information is now available about the properties of GDR built on the ground states and the excited states of nuclei revealing important properties of collective motion of nuclear many-body systems at zero and finite temperature [2, 3]. Understanding the structure of the nuclear collective response, its fragmentation and damping mechanism constitute a challenge for theoretical models [4–6]. To achieve this goal, one possible avenue is development of quantum transport models for nuclear collective dynamics [7, 8].

The theoretical investigations of nuclear collective response employing the random-phase approximation (RPA) based on the mean-field theory have been quite successful in describing mean resonance energies [9]. However, the RPA theory is not suitable for describing damping of collective excitations. There are different mechanisms involved in damping of nuclear collective states. A part of damping is due to particle emission giving rise to the escape width. The collective mode also acquires a spreading width as a consequence of the coupling with the internal degrees of freedom. In general, the spreading width is made up by three different contributions: *(i)* the Landau Damping due to spreading of the collective modes over non-collective particle–hole excitations, *(ii)* the coherent mechanism due to coupling with low-lying surface modes [4, 10], and *(iii)* the damping due to coupling with incoherent $2p$–$2h$ states usually referred to as the collisional damping [11,12].

Most investigations of the nuclear response carried out thus far are based on either the coherent damping mechanism or the collisional damping. The coherent mechanism is particularly important at low temperature and accounts for the main features of the collective response [10, 13]. On the other hand, the collisional damping is relatively weak at low temperature but its magnitude becomes larger with increasing temperature [11, 12]. In a recent work, nuclear collective response in different atomic nuclei was investigated on the basis of a stochastic transport theory, which incorporates both the coherent mechanism and the collisional damping in a consistent manner [14,15]. Such an extended quantum RPA calculations require a large amount of numerical effort and also may not provide a clear understanding of relative importance of different dissipation mechanisms. In this work, employing the same transport theory, we calculate the isovector response in nuclear matter at finite temperature and investigate the relative importance of the coherent damping mechanism and the collisional damping as a function of temperature.

Temporal evolution of the average value of the single-particle density matrix $\rho(t)$ in the stochastic transport theory is determined by a transport equation [16–18]

$$i\hbar\frac{\partial}{\partial t}\rho - [h(\rho), \rho] = K_{\mathrm{I}}(\rho) + K_{\mathrm{C}}(\rho)\,, \qquad (1)$$

where $h(\rho)$ is an effective mean-field Hamiltonian. $K_{\mathrm{I}}(\rho)$ represents the incoherent binary collision term, and $K_{\mathrm{C}}(\rho)$ is referred to as the coherent collision term. The coherent collision term describes the coupling between mean-field fluctuations and the single-particle motion. In the stochastic transport theory, this collision term has been analyzed by employing a time-dependent RPA approach, and shown that, it has the structure of particle–phonon collision term [17–19]. The linear response of the system to an external perturbation can be described by considering the small amplitude limit of the transport equation given in Eq. (1). The small deviations of the density matrix $\delta\rho(t) = \rho(t) - \rho_0$ around a finite temperature equilibrium state ρ_0 are determined by the linearized form of Eq. (1),

$$i\hbar\frac{\partial}{\partial t}\delta\rho - [h_0, \delta\rho] - [\delta U + F, \rho_0] = \delta K_{\mathrm{I}}(\rho) + \delta K_{\mathrm{C}}(\rho)\,. \qquad (2)$$

In this expression, $\delta U = (\partial U/\partial\rho)_0\,\delta\rho$ represents small deviations in the effective mean-field potential and $\delta K_{\mathrm{I}}(\rho)$ and $\delta K_{\mathrm{C}}(\rho)$ represents the linearized form of the non-Markovian incoherent and coherent collision terms, respectively. In order to study isovector collective response of the system, we include an external harmonic perturbation in the form $\tau_3 F[e^{-i\omega t} + e^{+i\omega t}]$, where τ_3 is the third component of the isospin operator.

In the isovector channel, small density fluctuations can be expressed as a difference in fluctuations of proton and neutron density matrices, $\delta\rho(t) = \delta\rho_p(t) - \delta\rho_n(t)$. In a similar manner, mean-field fluctuations in the isovector channel can be written as $\delta U = \delta U_p(\vec{r},t) - \delta U_n(\vec{r},t) = 2V_0\delta n(\vec{r},t)$, where $\delta n(\vec{r},t) = \delta n_p(\vec{r},t) - \delta n_n(\vec{r},t)$ denotes the local density fluctuations. We look for a solution of Eq. (1) of the form $\delta\rho(t) = \delta\rho(\omega)e^{-iwt} + \text{h.c.}$ In the small amplitude limit, collision terms are also harmonic, $\delta K_{\mathrm{I}} = \delta K_{\mathrm{I}}(\omega)e^{-iwt} + \text{h.c.}$ and $\delta K_{\mathrm{C}} = \delta K_{\mathrm{C}}(\omega)e^{-iwt} + \text{h.c.}$ In the homogeneous nuclear matter, the plane waves are the eigenmodes of the mean-field Hamiltonian h_0, and therefore provide a suitable representation to investigate the response of nuclear matter. In the momentum representation we then obtain [20, 21],

$$\begin{aligned}
&\left[\hbar\omega + \epsilon\left(\vec{p} - \frac{\vec{k}}{2}\right) - \epsilon\left(\vec{p} + \frac{\vec{k}}{2}\right)\right]\left\langle \vec{p} + \frac{\vec{k}}{2} \mid \delta\rho(w) \mid \vec{p} - \frac{\vec{k}}{2}\right\rangle \\
&- \left[f\left(\vec{p} - \frac{\vec{k}}{2}\right) - f\left(\vec{p} + \frac{\vec{k}}{2}\right)\right] \\
&\times \left\{\left\langle \vec{p} + \frac{\vec{k}}{2} \mid \delta U \mid \vec{p} - \frac{\vec{k}}{2}\right\rangle + 2\left\langle \vec{p} + \frac{\vec{k}}{2} \mid F(\vec{r}) \mid \vec{p} - \frac{\vec{k}}{2}\right\rangle\right\} \\
&= \left\langle \vec{p} + \frac{\vec{k}}{2} \mid \delta K_{\mathrm{I}}(\omega) \mid \vec{p} - \frac{\vec{k}}{2}\right\rangle + \left\langle \vec{p} + \frac{\vec{k}}{2} \mid \delta K_{\mathrm{C}}(\omega) \mid \vec{p} - \frac{\vec{k}}{2}\right\rangle, \qquad (3)
\end{aligned}$$

where $f(\vec{q}) = 1/\left\{1 + e^{\beta[\epsilon(\vec{q}) - \mu]}\right\}$ is the Fermi–Dirac occupation factor. The isovector response function can be deduced from Eq. (3) by integrating over the momentum $\vec{p}$ and using the fact that the Fourier transform of the local isovector density fluctuations is related to the fluctuations of the density matrix according to,

$$\int 2\frac{d^3p}{(2\pi\hbar)^3}\left\langle \vec{p}+\frac{\vec{k}}{2}\mid \delta\rho(w)\mid \vec{p}-\frac{\vec{k}}{2}\right\rangle = \delta n(\vec{k},w)\,. \tag{4}$$

The integration over the momentum, then, yields

$$\begin{aligned} &\delta n(\vec{k},w) - \left[V_0\delta n(\vec{k},w)+F(\vec{k})\right] \mathit{\Pi}_0(\vec{k},w) \\ &= \left[V_0\delta n(\vec{k},w)+F(\vec{k})\right] \mathit{\Pi}_{\rm I}(\vec{k},w) + \left[V_0\delta n(\vec{k},w)+F(\vec{k})\right] \mathit{\Pi}_{\rm C}(\vec{k},w)\,, \end{aligned} \tag{5}$$

where $F(\vec{k})$ denotes the Fourier transform of the external perturbation. In this expression the quantity $\mathit{\Pi}_0(\vec{k},w)$ is the unperturbed Lindhard function,

$$\mathit{\Pi}_0(\vec{k},w) = \frac{4}{(2\pi\hbar)^3}\int d^3p\, \frac{f\left(\vec{p}-\frac{\vec{k}}{2}\right)-f\left(\vec{p}+\frac{\vec{k}}{2}\right)}{\hbar w - \epsilon\left(\vec{p}+\frac{\vec{k}}{2}\right)+\epsilon\left(\vec{p}-\frac{\vec{k}}{2}\right)+i\gamma}\,. \tag{6}$$

The other quantities, $\mathit{\Pi}_{\rm I}(\vec{k},w)$ and $\mathit{\Pi}_{\rm C}(\vec{k},w)$ are the collisional and coherent response functions, respectively, and they are related to the incoherent and coherent collision terms according to,

$$\begin{aligned} &\left[V_0\delta n(\vec{k},w)+F(\vec{k})\right] \mathit{\Pi}_{\rm I}(\vec{k},w) \\ &= \frac{2}{(2\pi\hbar)^3}\int\, d^3p\, \frac{\left\langle \vec{p}+\frac{\vec{k}}{2}\mid \delta K_{\rm I}(w)\mid \vec{p}-\frac{\vec{k}}{2}\right\rangle}{\hbar w - \epsilon\left(\vec{p}+\frac{\vec{k}}{2}\right)+\epsilon\left(\vec{p}-\frac{\vec{k}}{2}\right)+i\gamma}\,, \end{aligned} \tag{7}$$

and

$$\begin{aligned} &\left[V_0\delta n(\vec{k},w)+F(\vec{k})\right] \mathit{\Pi}_{\rm C}(\vec{k},w) \\ &= \frac{2}{(2\pi\hbar)^3}\int\, d^3p\, \frac{\left\langle \vec{p}+\frac{\vec{k}}{2}\mid \delta K_{\rm C}(w)\mid \vec{p}-\frac{\vec{k}}{2}\right\rangle}{\hbar w - \epsilon\left(\vec{p}+\frac{\vec{k}}{2}\right)+\epsilon\left(\vec{p}-\frac{\vec{k}}{2}\right)+i\gamma}\,. \end{aligned} \tag{8}$$

The retarded response function, which is defined by

$$\delta n(\vec{k},w) = \mathit{\Pi}_{\rm R}(\vec{k},w)F(\vec{k}) \tag{9}$$

is then obtained as

$$\mathit{\Pi}_{\rm R}(\vec{k},w) = \frac{\mathit{\Pi}(\vec{k},w)}{1-V_0\mathit{\Pi}(\vec{k},w)}\,, \tag{10}$$

where $\Pi(\vec{k},w) = \Pi_0(\vec{k},w) + \Pi_{\rm I}(\vec{k},w) + \Pi_{\rm C}(\vec{k},w)$. The strength distribution function is given by the imaginary part of the retarded response function as,

$$S(\vec{k},w) = -\frac{1}{\pi}{\rm Im}\Pi_{\rm R}(\vec{k},w)\,. \tag{11}$$

In a previous work, we derived an explicit expression for the collisional response function $\Pi_{\rm I}(\vec{k},w)$ [11,12,21] and investigated the collisional damping on the giant dipole excitations. The collisional response function is given by,

$$\begin{aligned}\Pi_{\rm I}(\vec{k},w) = & \frac{2}{(2\pi\hbar)^3}\int d^3p_1 d^3p_2 d^3p_3 d^3p_4 \left|\frac{\Delta Q}{2}\right|^2 \frac{W(12;34)}{\pi} \\ & \times \frac{f_1 f_2 \overline{f}_3 \overline{f}_4 - \overline{f}_1 \overline{f}_2 f_3 f_4}{\hbar w - \epsilon_3 - \epsilon_4 + \epsilon_1 + \epsilon_2 + i\eta}\,,\end{aligned} \tag{12}$$

where $\Delta Q = Q_1 + Q_2 - Q_3 - Q_4$ with $Q_i = 1/\left[\hbar w - \epsilon\left(\vec{p}_i + \frac{\vec{k}}{2}\right) + \epsilon\left(\vec{p}_i - \frac{\vec{k}}{2}\right)\right]$, $\overline{f}_i = 1 - f_i$ and $W(12;34)$ denotes the basic two-body transition rate, which can be expressed in terms of the spin averaged proton–neutron scattering cross section as

$$W(12;34) = \frac{1}{(2\pi\hbar)^3}\,\frac{4\hbar}{m^2}\,\left(\frac{d\sigma}{d\Omega}\right)_{pn}\,\delta^3(\vec{p}_1 + \vec{p}_2 - \vec{p}_3 - \vec{p}_4)\,. \tag{13}$$

As stated above, the stochastic transport theory involves, in addition to the collisional damping, a coherent dissipation mechanism due to coupling mean-field fluctuations. In the special case of small amplitude density fluctuations, the coherent dissipation mechanism takes the form of particle–phonon collision term. In references [17,18], in terms of a time-dependent RPA formalism an explicit expression of the coherent collision term was presented. In Appendix A, we present a brief description of the linearized form of the coherent collision term. Following from the Eq. (26), the coherent response function in nuclear matter is given by

$$\begin{aligned}\Pi_{\rm C}(\vec{k},w) = & \frac{4}{2\pi\hbar}\frac{1}{(2\pi\hbar)^6}\int d^3p_2\; d^3p_1\; dW \left|\frac{\partial U}{\partial n}\right|_0^2 \tilde{\sigma}(\vec{q},W)|Q_2 - Q_1|^2 \\ & \times\left[\frac{(N_W{+}1)\left\{f\left(\vec{p}_2 - \frac{\vec{k}}{2}\right)\left[1 - f\left(\vec{p}_1 + \frac{\vec{k}}{2}\right)\right]\right\} - N_W\left\{f\left(\vec{p}_1 + \frac{\vec{k}}{2}\right)\left[1 - f\left(\vec{p}_2 - \frac{\vec{k}}{2}\right)\right]\right\}}{\hbar W - \hbar\omega + \epsilon\left(\vec{p}_1 + \frac{\vec{k}}{2}\right) - \epsilon\left(\vec{p}_2 - \frac{\vec{k}}{2}\right) + i\eta}\right],\end{aligned} \tag{14}$$

where $\vec{q} = \vec{p}_2 - \vec{p}_1$ and $N_W = 1/[e^{\hbar W/T} - 1]$ is the phonon occupation number. In this expression $U = (U_p + U_n)/2$ represents the mean-field potential as a function of the total nucleon density $n = n_p + n_n$, and its derivative is evaluated at the equilibrium density n_0. The quantity $\tilde{\sigma}(\vec{q}, W)$ denotes the isoscalar density correlation function. It is possible to calculate the density correlation function in the framework of the stochastic transport model. Here, we give the result obtained in the mean-field approximation, and for details we refer [19],

$$\tilde{\sigma}(\vec{q}, W) = -\frac{2\mathrm{Im}[\Pi_0(\vec{q}, W)]}{\left[1 - \left(\frac{\partial U}{\partial n}\right)_0 Re[\Pi_0(\vec{q}, W)]\right]^2 + \left[\left(\frac{\partial U}{\partial n}\right)_0 \mathrm{Im}[\Pi_0(\vec{q}, W)]\right]^2} . \quad (15)$$

In the calculations, we employ a simplified Skyrme interaction [4, 22],

$$v = t_0(1 + x_0 P_\sigma)\delta(\vec{r}) + \frac{1}{6} t_3(1 + x_3 P_\sigma)\rho^\alpha(\vec{R})\delta(\vec{r}) , \quad (16)$$

with $\vec{r} = \vec{r}_1 - \vec{r}_2$ and $\vec{R} = (\vec{r}_1 + \vec{r}_2)/2$. The local potential for protons is then given by

$$\begin{aligned} U_p(\vec{r}, t) = \quad & t_0 \left(1 + \frac{1}{2} x_0\right) n(\vec{r}, t) - t_0 \left(\frac{1}{2} + x_0\right) n_p(\vec{r}, t) \\ & + \frac{1}{12} t_3 n^\alpha(\vec{r}, t) \left[(2 + \alpha)\left(1 + \frac{1}{2} x_3\right) n(\vec{r}, t) \right. \\ & \left. -2\left(\frac{1}{2} + x_3\right) n_p(\vec{r}, t) - \alpha \left(\frac{1}{2} + x_3\right) \frac{n(\vec{r}, t)}{2}\right] , \end{aligned} \quad (17)$$

with a similar expression for neutrons. Then the derivative of the mean-field potential quantity for symmetric nuclear matter in equilibrium is given by,

$$\left(\frac{\partial U}{\partial n}\right)_{n=n_0} = \frac{3t_0}{4} + \frac{1}{16} t_3 n_0^\alpha (\alpha + 1)(\alpha + 2) . \quad (18)$$

In the linear response approximation, the coupling constant for isovector dipole vibrations becomes,

$$V_0 = -\frac{1}{2} t_0 \left(\frac{1}{2} + x_0\right) - \frac{1}{12} t_3 n_0^\alpha \left(\frac{1}{2} + x_3\right) . \quad (19)$$

In our analysis, we consider two different effective Skyrme interactions: SLy4 and SV forces. The parameters for SLy4 are $t_0 = -2488.91$ MeV fm^3, $t_3 = 13777$ MeV fm$^{7/2}$, $x_0 = 0.834$, $x_3 = 1.354$ and $\alpha = 1/6$ which results in the value $V_0 = 85$ MeV [20], and the corresponding parameters for SV are $t_0 = -1248.29$ MeV fm^3, $t_3 = 0.0$ MeV fm$^{7/2}$, $x_0 = -0.17$, $x_3 = 1.0$ and $\alpha = 1$ which results in the value $V_0 = 205.97$ MeV [9].

As a result of the approximate treatment, the response functions $\Pi_{\rm I}(\vec{k},\omega)$ and $\Pi_{\rm C}(\vec{k},\omega)$ have singular behavior arising from the pole of the distortion functions,

$$Q_i = 1 \Bigg/ \left[\hbar w - \epsilon\left(\vec{p_i} + \frac{\vec{k}}{2}\right) + \epsilon\left(\vec{p_i} - \frac{\vec{k}}{2}\right)\right] .$$

In the previous work [12, 21], this singular behavior was avoided by incorporating a pole approximation. Here consider a similar approximation. In the distortion functions, we make the replacement $\omega \to \omega_{\rm D} - i\Gamma/2$ where $\omega_{\rm D}$ and Γ are the mean frequency and the width of the resonance, respectively. In the previous work these quantities are determined by solving the dispersion relation $1 - V_0\Pi_1(\vec{k},\omega) = 0$ at each temperature that is considered. In this work, rather than solving the dispersion relation, we take the mean frequency $\omega_{\rm D}$ and the width Γ directly from the RPA strength functions. Furthermore, we neglect the real parts of response functions $\Pi_{\rm I}(\vec{q},W)$ and $\Pi_{\rm C}(\vec{q},W)$ in our calculations. In the calculations of the collisional response function $\Pi_{\rm I}(\vec{q},W)$, using conservation laws and symmetry properties, it is possible to reduce the twelve dimensional integrals to five fold integrals by incorporating the transformations into the total momenta $\vec{P} = \vec{p}_1 + \vec{p}_2$, $\vec{P'} = \vec{p}_3 + \vec{p}_4$, and relative momenta $\vec{q} = (\vec{p}_1 - \vec{p}_2)/2$, $\vec{q}\,' = (\vec{p}_3 - \vec{p}_4)/2$ before and after the collisions. The integral over $\vec{P'}$ can be performed immediately. The delta function $\delta(\hbar\omega - \epsilon' + \epsilon)$ in ${\rm Im}\Pi_{\rm I}(\vec{k},w)$ where $\epsilon = \vec{q}\,^2/m$ and $\epsilon' = \vec{q}\,'\,^2/m$ are the energies of two particle system in the center of mass frame before and after the collision makes it possible the reduce the integrals further using familiar methods from the Fermi liquid theory [22]. Then, we evaluate the remaining five dimensional integrals numerically by employing a fast algorithm. In the evaluation of momentum integrals, we neglect the angular anisotropy of the cross sections and make the replacement $(d\sigma/d\Omega)_{pn} \to \sigma_{pn}/4\pi$ with $\sigma_{pn} = 40$ mb. In the calculations of the imaginary part of the coherent response function ${\rm Im}\Pi_{\rm C}(\vec{q},W)$, it is convenient to choose the z-axis along direction of the wave vector $\vec{k}$ and carry out the momentum integrals in spherical representation, where $d^3p_1 = \sqrt{2m\epsilon_1}m\epsilon_1 \sin\theta_1 d\theta_1 d\phi_1$ and $d^3p_2 = \sqrt{2m\epsilon_2}m\epsilon_2 \sin\theta_2 d\theta_2 d\phi_2$. Similar to the collisional response, in the coherent response function at low temperatures, particle–hole excitations are concentrated around the Fermi surface. Also, for the magnitude of the wave vector that is much smaller than the Fermi momentum, $k \ll p_{\rm F}$, except in distortion functions Q, we can ignore the $\vec{k}$ dependence in Eq. (14), and carry out the energy integrals over ϵ_1 and ϵ_2, analytically. Remaining four integrals over the angles θ_1, θ_2 and $\phi_2 - \phi_1$, and over the frequency W are done numerically.

In order to apply our results to finite nuclei, we work within the framework of Steinwedel and Jensen model for nuclear dipole oscillations [9, 23]. In this model neutrons and protons oscillate inside a sphere of radius R given by the expression

$$\rho_p(\vec{r},t) - \rho_n(\vec{r},t) = F \sin(\vec{k}\cdot\vec{r}) e^{iwt}\,, \tag{20}$$

the total density remaining equal to the saturation density ρ_0 of nuclear matter and the wavenumber k is given by $k = \pi/2R$. We apply Steinwedel and Jensen model to GDR in ^{120}Sn and ^{208}Pb, and we take $R = 5.6$ fm $k = 0.28$ fm^{-1} for ^{120}Sn and $R = 6.7$ fm $k = 0.23$ fm^{-1} for ^{208}Pb according to $R = 1.13A^{1/3}$. We show our results for the strength function with and without the collision terms in Figs. 1 and 3 for ^{120}Sn, and in Figs. 2 and 4 for ^{208}Pb as a function of experimental temperature T^* where we also present the comparison with the normalized experimental data [3]. The experimental temperature T^* is related to the temperature parameter in the Fermi–Dirac function $f(\epsilon, T)$ as $T = T^*\sqrt{a_{\rm E}/a_{\rm F}}$, where $a_{\rm E}$ denotes the energy dependent empirical level density parameter and $a_{\rm F} = A\pi^2/4\epsilon_{\rm F}$ denotes Fermi gas level density parameter [3, 24]. In our calculations, we use the temperature values T in the Fermi–Dirac function that are related to the experimental temperatures in this manner.

We observe that the result of calculations are rather sensitive to the effective interactions, even in the RPA level. At the RPA level, without the contribution of the collision terms, the calculated peak position of the strength functions do not change with temperature. As a matter of fact, as seen from Figs. 1 and 2, calculations with SLy4 force we find for ^{208}Pb the peak is at $\omega = 12.3$ MeV for $T^* = 1.34,\ 1.62,\ 1.85,\ 2.05$ MeV while for ^{120}Sn it occurs at $\omega = 14.3$ MeV for $T^* = 1.36,\ 2.13$ MeV and at $\omega = 14.2$ MeV for $T^* = 2.67,\ 3.1$ MeV. These results are consistent with the calculations reported in our previous work [21]. As seen in Figs. 3 and 4, we find a similar behavior in calculations with SV force, however we observe that peak values are shifted to somewhat higher values, $\omega = 13.0$ MeV for ^{208}Pb and $\omega = 15.0$ MeV for ^{120}Sn. This behavior of the peak energy with temperature is in accordance with the experimental results, where it is observed that the mean-energy of the dipole response is almost constant for ^{208}Pb when T^* changes between 1.3 to 2.0 MeV while a decrease of 1.5 MeV is observed in ^{120}Sn when T^* changes from 1.2 to 3.1 MeV [3]. Position of the peak value is rather sensitive to the effective force employed. In particular with SLy4 force, the average positions of the peak values of the strength functions are slightly below the experimental values, which are $\omega = 15.4$ MeV for ^{120}Sn and $\omega = 13.4$ MeV for ^{208}Pb. Moreover, the value for $k = \pi/2R$ that is used in Steinwedel and Jensen model depends on the value of R_0 used in $R = R_0A^{1/3}$, and changing k somewhat also produces a change in the position of the peak but in general the above conclusions are not affected. Furthermore, these results for the peak position of the strength functions are in accordance with the earlier RPA calculations [9].

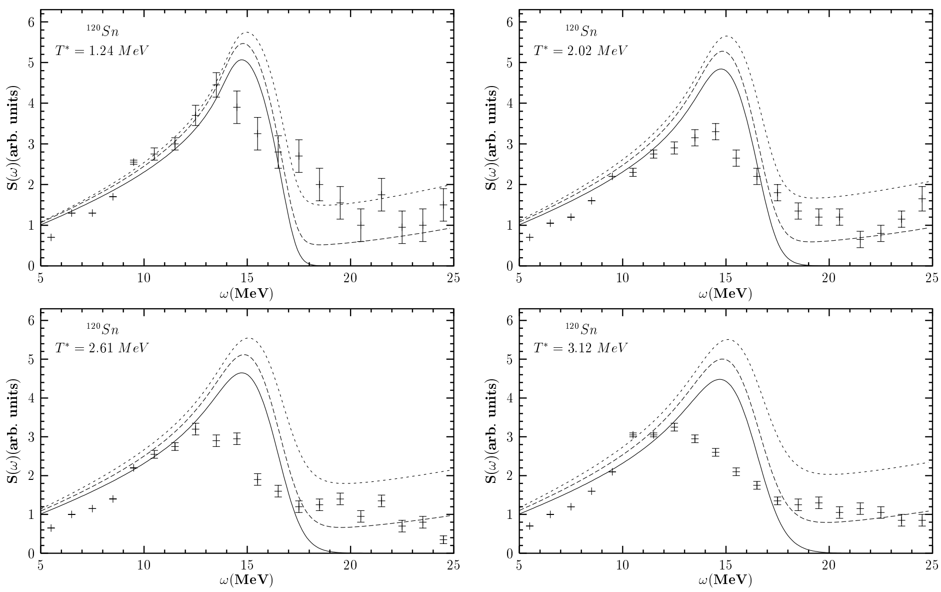

Fig. 1. The GDR strength function of ^{120}Sn obtained using Steinwedel–Jensen model at different temperatures. Solid-, dashed- and dotted-lines show the results of the calculations with the effective SLy4 Skyrme interaction in the mean-field approximation, including the collisional damping and including both collisional and coherent damping, respectively.

Dashed-lines in figures show the strength functions including collisional damping mechanism. Since, we neglect the real part of the collisional response the peak values of the strength do not change, but the collisional mechanism introduces a small spread in the strength distributions at low frequencies. At high frequency tail of the strength functions, the collisional damping appears to be stronger, however its magnitude is weaker than the results reported in the previous work [21]. This difference arises from a different manner of the implementation of the pole approximation, as explained above. Dotted-lines in Figs. 1 and 3, and Figs. 3 and 4, show the strength functions including both the collisional and coherent damping mechanisms for ^{120}Sn and ^{208}Pb, respectively. As expected, calculations indicate that the coherent mechanism due to coupling with isoscalar sound modes have a stronger influence on the damping of the isovector vibrations.

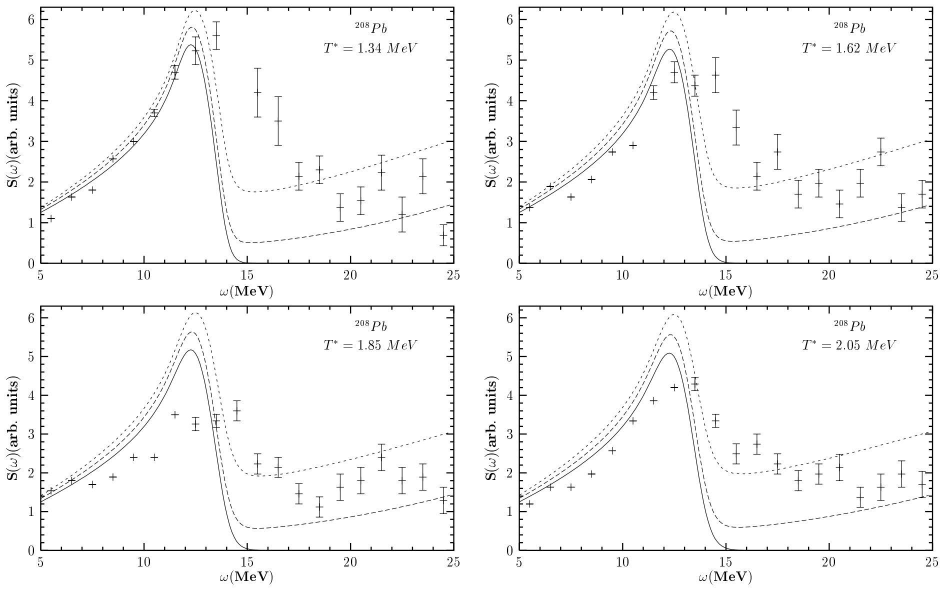

Fig. 2. The GDR strength function of ^{208}Pb obtained using Steinwedel–Jensen model at different temperatures. Solid-, dashed- and dotted-lines show the results of the calculations with the effective SLy4 Skyrme interaction in the mean-field approximation, including the collisional damping and including both collisional and coherent damping, respectively.

Rather simple description presented in this paper, including collisional and coherent dissipation mechanism in nuclear matter, is able to explain certain qualitative properties of giant dipole excitations in ^{120}Sn and ^{208}Pb at finite temperature. The result of calculations are rather sensitive the effective force employed, and it appears that SLy4 force seems to provide a better description of the giant dipole strength functions. However, calculations do not produce a quantitative description of the experimental strength functions as a function of temperature, which are indicated by bars in figures. We also note that the visible effect of the total strength non-conservation in Figs. 1–4, when compared the results in the mean-field approximation, then including the collisional damping and then in addition the coherent damping. This sum rule non-conservation is partly due to omission of the real part in the propagator, since this leads to modification of underlying dispersion relation. Including the real part of the propagator and using Skyrme interaction within the present formalism, one would have to refit the interaction parameters to reproduce the experimental centroid energies, which is not done in the present work. This sum rule non-conservation is partly due to pole approximation, which is introduced to avoid the singular

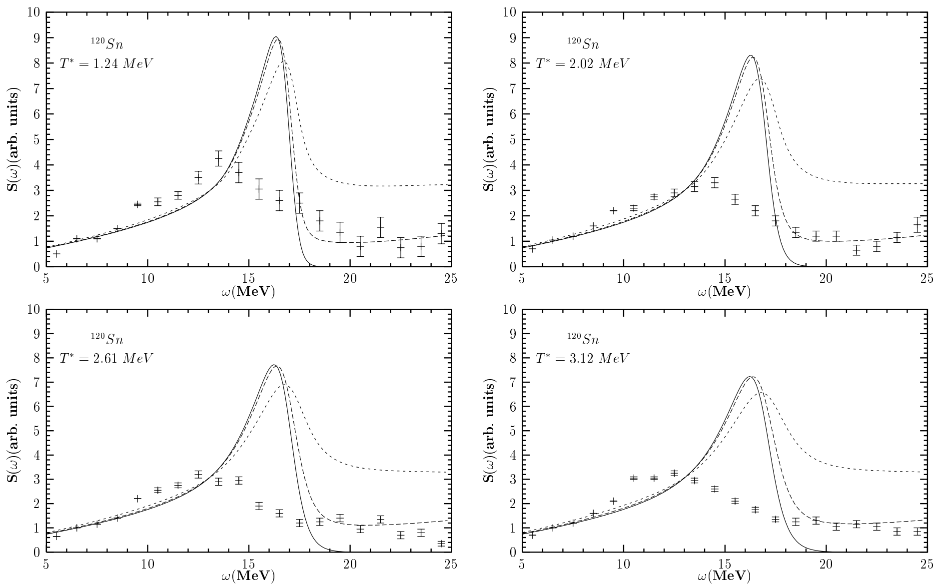

Fig. 3. The GDR strength function of ^{120}Sn obtained using Steinwedel–Jensen model at different temperatures. Solid-, dashed- and dotted-lines show the results of the calculations with the effective SV Skyrme interaction in the mean-field approximation, including the collisional damping and including both collisional and coherent damping, respectively.

behavior of the distortion functions in the collisional and as well as coherent response functions. In particular, rapid rise of the strength functions at high frequencies is originating from this approximation. Another important element is missing in the calculations is the surface effect in the coherent damping mechanism. It will be interesting to improve the description by incorporating the real part in the propagator in the response functions, and also by including coupling of dipole vibrations with low frequency surface modes which maybe carried out in the Thomas–Fermi framework [4, 10, 15].

One of us (S.A.) gratefully acknowledges TUBITAK for support and the Physics Department of the Middle East Technological University for warm hospitality extended to him during his visit. This work is supported in part by the US DOE grant No. DE-FG05-89ER40530.

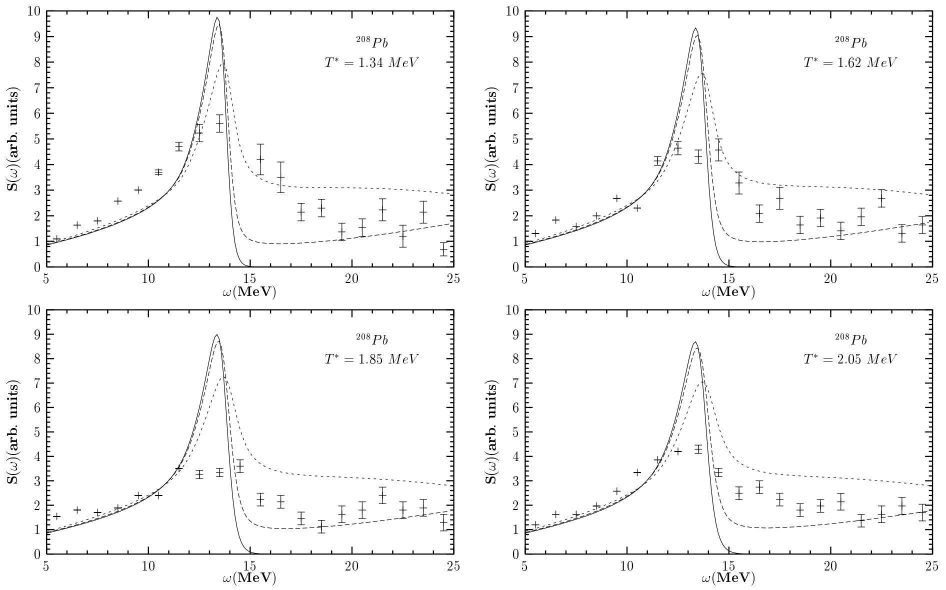

Fig. 4. The GDR strength function of ^{208}Pb obtained using Steinwedel–Jensen model at different temperatures. Solid-, dashed- and dotted-lines show the results of the calculations with the effective SV Skyrme interaction in the mean-field approximation, including the collisional damping and including both collisional and coherent damping, respectively.

Appendix A

In the stochastic transport theory, the coupling between the mean fluctuations and the single-particle motion gives rise to a dissipation mechanism and the associated collision terms takes the form of a particle–phonon collision term. This coherent collision term has been investigated before in quantal and semi-classical framework. In the linear response regime, it is sufficient to consider the linearized form of the coherent collision term around equilibrium, which can be expressed as

$$\delta K_{\mathrm{C}} = \int_0^\infty dt' \, \overline{\left[\delta\hat{U}(t), [G_0(t,t')[\delta\Phi(t'), \delta\hat{U}(t')]G_0(t,t'), \rho_0]\right]} \, . \qquad (21)$$

In this expression, the bar indicates an ensemble averaging, $\delta\hat{U}(t) = (\partial U/\partial n)_0 \delta\hat{n}(t)$ denotes small amplitude isoscalar fluctuations of the mean-field around its equilibrium value and $G_0 = \exp[-ih_0(t-t')]$ represents the mean-field propagator. The quantity $\delta\Phi(t)$ is the RPA amplitude associated with the harmonic density vibrations, and it is related to the harmonic vibrations of the density matrix according to $\delta\rho(t) = [\delta\Phi(t), \rho_0]$. Here we consider harmonic isovector vibrations in a homogenous nuclear matter. In the momentum representation, with the help of Eq. (3), we can determine the matrix elements of the isovector RPA amplitudes to give,

$$\left\langle \vec{p}+\frac{\vec{k}}{2} \mid \delta\Phi(t) \mid \vec{p}-\frac{\vec{k}}{2}\right\rangle = 2e^{-i\omega t}\left[\frac{V_0\delta n(\vec{k},\omega)+F(\vec{k})}{\hbar\omega-\epsilon\left(\vec{p}+\frac{\vec{k}}{2}\right)+\epsilon\left(\vec{p}-\frac{\vec{k}}{2}\right)+i\eta}\right]. \quad (22)$$

In obtaining this result we neglect the damping terms on the right hand side of Eq. (3).

In order to calculate the variance of small amplitude density fluctuations, it is convenient to introduce time Fourier transform of the mean-field fluctuations,

$$\left\langle \vec{p}+\frac{\vec{k}}{2} \mid \delta\hat{U}(t) \mid \vec{p}-\frac{\vec{k}}{2}\right\rangle = \int_{-\infty}^{\infty}\frac{d\omega}{2\pi}\left(\frac{\partial U}{\partial n}\right)_0 \delta\hat{n}(\vec{k},\omega)\, e^{-i\omega t}. \quad (23)$$

In the framework of the stochastic transport theory, it is possible to calculate the variance of the density fluctuations to give

$$\overline{\delta\hat{n}(\vec{k}_1,\omega_1)\delta\hat{n}(\vec{k}_2,\omega_2)} = (2\pi\hbar)^4\delta(\vec{k}_1+\vec{k}_2)\delta(\omega_1+\omega_2)\tilde{\sigma}(\vec{p}_1,\omega_1)\left(N_{\omega_1}+\frac{1}{2}\right), \quad (24)$$

where the density correlation function $\sigma(\vec{p}_1,\omega_1)$ is given by Eq. (15) and N_{ω_1} denotes the phonon occupation number with the frequency ω_1. Using this result for the density correlation function and inserting appropriate intermediate momentum states into Eq. (22), the matrix elements of the coherent collision term is expressed as

$$\begin{aligned}
\left\langle \vec{p}_2+\frac{\vec{k}}{2} \mid \delta K_{\mathrm{C}}(w) \mid \vec{p}_2-\frac{\vec{k}}{2}\right\rangle &= 2\int\frac{d^3p_1}{(2\pi\hbar)^4}dW\left|\frac{\partial U}{\partial n}\right|^2[\tilde{\sigma}(\vec{p}_2-\vec{p}_1,W)] \\
&\times\left[\frac{\left(N_W+\frac{1}{2}\right)\left[f\left(\vec{p}_2-\frac{\vec{k}}{2}\right)-f\left(\vec{p}_1+\frac{\vec{k}}{2}\right)\right]}{\hbar W-\hbar\omega+\epsilon\left(\vec{p}_1+\frac{\vec{k}}{2}\right)-\epsilon\left(\vec{p}_2-\frac{\vec{k}}{2}\right)+i\eta}\right] \\
&\times(Q_1-Q_2)(V_0\delta n(\vec{k},\omega)+F(\vec{k},\omega)) \\
&+2\int\frac{d^3p_1}{(2\pi\hbar)^4}dW\left|\frac{\partial U}{\partial n}\right|^2\tilde{\sigma}(\vec{p}_1-\vec{p}_2,W) \\
&\times\left[\frac{\left(N_W+\frac{1}{2}\right)\left[f\left(\vec{p}_2+\frac{\vec{k}}{2}\right)-f\left(\vec{p}_1-\frac{\vec{k}}{2}\right)\right]}{\hbar W-\hbar\omega+\epsilon\left(\vec{p}_2+\frac{\vec{k}}{2}\right)-\epsilon\left(\vec{p}_1-\frac{\vec{k}}{2}\right)+i\eta}\right] \\
&\times(Q_1-Q_2)(V_0\delta n(\vec{k},\omega)+F(\vec{k},\omega)),
\end{aligned} \quad (25)$$

where $Q_i = 1/\hbar\omega - \epsilon\left(\vec{p}_i + \frac{\vec{k}}{2}\right) + \epsilon\left(\vec{p}_i - \frac{\vec{k}}{2}\right) + i\gamma$. Using this expression the coherent response function $\Pi_{\mathrm{C}}(\vec{k}, w)$ can be expressed as,

$$\Pi_{\mathrm{C}}(\vec{k}, w) = \frac{4}{2\pi\hbar}\frac{1}{(2\pi\hbar)^6}\int d^3p_2\, d^3p_1\, dW \left|\frac{\partial U}{\partial n}\right|^2 \tilde{\sigma}(\vec{q}, W)\,|Q_2 - Q_1|^2$$
$$\times\left[\frac{(N_W+\frac{1}{2})\left\{f\left(\vec{p}_2-\frac{\vec{k}}{2}\right)\left[1-f(\vec{p}_1+\frac{\vec{k}}{2})\right]\right\}-(N_W+\frac{1}{2})\left\{f\left(\vec{p}_1+\frac{\vec{k}}{2}\right)\left[1-f\left(\vec{p}_2-\frac{\vec{k}}{2}\right)\right]\right\}}{\hbar W-\hbar\omega+\epsilon\left(\vec{p}_1+\frac{\vec{k}}{2}\right)-\epsilon\left(\vec{p}_2-\frac{\vec{k}}{2}\right)+i\eta}\right], \tag{26}$$

where $\vec{q} = \vec{q}_2 - \vec{q}_1$, and the Fermi–Dirac factors are written in a symmetrical form. By inspection, it can be seen that in this expression, the first and second terms correspond to excitation and absorption of sound phonons. These rates should be proportional to $N_W + 1$ and N_W, respectively, but the average value $N_W + 1/2$ appears in both rates. There are other contributions in the coherent collision term $K_{\mathrm{C}}(\rho)$ in Eq. (1) arising from the cross correlations between the collective and non collective modes. In schematic models, it is possible to show that these cross correlations give rise to additional contributions to the collision term, so that the excitation and absorption rates become proportional to $N_W + 1$ and N_W, as it should be [17, 18]. However, in the RPA analysis , it is difficult to extract such contributions. For the time being, we replace the excitation and absorbtion factors by $N_W + 1$ and N_W, and express the coherent respond function as given by Eq. (15).

REFERENCES

[1] *Electric and Magnetic Giant Resonances in Nuclei*, J. Speth ed., World Scientific, Singapore 1991.

[2] E. Ramakrishan *et al.*, *Phys. Rev. Lett.* **76**, 2025 (1996); *Phys. Lett.* **B383**, 252 (1996).

[3] T. Baumann *et al.*, *Nucl. Phys.* **A635**, 428 (1998).

[4] G.F. Bertsch, P.F. Bortignon, R.A. Broglia, *Rev. Mod. Phys.* **55**, 287 (1983).

[5] G.F. Bertsch, R.A. Broglia, *Oscillations in Finite Quantum Systems*, Cambridge 1994.

[6] S. Drozdz, S. Nishizaki, J. Speth, J. Wambach, *Phys. Rep.* **197**, 1 (1990).

[7] K. Goeke, P.G. Reinhard, Time-Dependent Hartree-Fock and Beyond, proceedings, Bad Honnef, Germany 1982.

[8] Y. Abe, S. Ayik, P.G. Reinhard, E. Suraud, *Phys. Rep.* **275**, 49 (1996).

[9] F.L. Braghin, D. Vautherin, A. Abada, *Phys. Rev.* **C52**, 2504 (1995).

[10] P.F. Bortignon, R.A. Bloglia, G.F. Bertsch, J. Pacheco, *Nucl. Phys.* **A460**, 149 (1986).

[11] S. Ayik, O. Yilmaz, A. Gokalp, P. Schuck, *Phys. Rev.* **C58**, 1594 (1998).

[12] O. Yilmaz, A. Gokalp, S. Yildirim, S. Ayik, *Phys. Lett.* **B472**, 258 (2000).

[13] S. Kamerdzhiev, J. Speth, G. Tertychny, *Nucl. Phys.* **A624**, 328 (1997).

[14] D. Lacroix, S. Ayik, Ph. Chomaz, *Phys. Rev.* **C63**, 064305 (2001).

[15] D. Lacroix, S. Ayik, Ph. Chomaz, *Prog. Part. Nucl. Phys.* **52**, 497 (2004).

[16] Y. Abe, S. Ayik, P.G. Reinhard, E. Suraud, *Phys. Rep.* **275**, 49 (1996).

[17] S. Ayik, *Phys. Lett.* **B493**, 47 (2000).

[18] S. Ayik, Y. Abe, *Phys. Rev.* **C64**, 024609 (2001).

[19] S. Ayik, *Z. Phys.* **A350**, 45 (1994).

[20] E. Chabanat, P. Bonche, P. Haensel, J. Meyer, R. Schaeffer, *Nucl. Phys.* **A627**, 710 (1997); *Nucl. Phys.* **A635**, 231 (1998).

[21] S. Ayik, A. Gokalp, O. Yilmaz, K. Bozkurt, *Acta Phys. Pol.* **B34**, 4229 (2003).

[22] G. Baym, C. Pethick, *Landau Fermi-Liquid Theory*, John Wiley, New York 1991.

[23] P. Ring, P. Schuck, *The Nuclear Many Body Problem*, Springer, Berlin 1980.

[24] D. Kusnezov *et al.*, *Phys. Rev. Lett.* **81**, 542 (1998).