Isovector Collective Response Function of Nuclear Matter at
Finite Temperature

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Abstract

We study isovector collective excitations in nuclear matter by employing the linearized Landau-Vlasov equation with and without a non-Markovian binary collision term at finite temperature. We calculate the giant dipole resonance (GDR) strength function for finite nuclei using Steinwedel-Jensen model and also by Thomas-Fermi approximation, and we compare them for $^{120}$Sn and $^{208}$Pb with experimental results.
Giant resonances, in particular giant dipole resonance (GDR), built on highly excited nuclear states have been the subject of many experimental and theoretical studies in recent years [1–3]. The damping properties and the excitation energy dependence of GDR width are still among the open problems in nuclear collective dynamics. There are essentially two different theoretical approaches to this problem. The first one explains the temperature dependence of the width by the coherent mechanism due to adiabatic coupling of the collective state with thermal surface deformations [4]. In the second approach, referred to as collisional damping, the coupling with incoherent two particle-two hole states plus a Landau damping form the mechanism for temperature dependence [5–11].

Investigation of the GDR strength function has been carried out in the mean-field approximation without the collisional damping in [10]. In a previous work, we investigated the GDR strength function in infinite nuclear matter in quantal framework by employing the linearized version of the extended time dependent Hartree-Fock theory (TDHF) with a non-Markovian binary collision term [12]. We then applied our results to finite nuclei using the Steinwedel-Jensen model. However, this model is not very reliable for the treatment of collective dipole oscillations of finite nuclei since in this model it is assumed that the nuclear surface remains constant and the density oscillations obeys a wave equation with the boundary condition that the radial velocity vanishes on the spherical surface with radius \( R = R_0 A^{1/3} \). On the other hand, semi-classical approaches based on Thomas-Fermi approximation which are easier to handle than quantal calculations have been very useful to tackle the problems related to finite nuclei [13]. Indeed, the long wave length limit of TDHF equation is the Landau-Vlasov equation which is valid if the spatial variations are slow. Moreover, if Thomas-Fermi method is used to determine the nuclear density then the Landau-Vlasov equation forms a better approximation to quantal theory even if the spatial variations are not small [13].

In this work, we study the GDR strength function of finite nuclei by employing the Landau-Vlasov equation. We do not consider the coherent mechanism and investigate the temperature dependence due to coupling of the collective state with incoherent two particle-
two hole states using a non Markovian collision term in the linearized Landau-Vlasov equation. We apply our results to finite nuclei using Steinwedel-Jensen model. We then consider giant dipole excitations in finite nuclei by expanding our results using Thomas-Fermi approximation. This way, we not only asses the effects of the collision term but we also compare the differences between infinite nuclear matter and finite nuclei results.

The transport equation [14], for the particle phase-space density $f(\vec{r}, \vec{p}, t)$ with collision term is the Boltzmann equation

$$\frac{\partial}{\partial t} f(\vec{r}, \vec{p}, t) + \vec{v}_p \epsilon(\vec{r}, \vec{p}, t). \vec{\nabla}_r f(\vec{r}, \vec{p}, t) - \vec{\nabla}_r \epsilon(\vec{r}, \vec{p}, t). \vec{\nabla}_p f(\vec{r}, \vec{p}, t) = K(f) \ .$$  \hspace{1cm} (1)

In the framework of Fermi liquid theory defining the quasiparticle velocity as $\vec{v} = \vec{\nabla}_p \epsilon(\vec{r}, \vec{p}, t)$ and assuming that the potential energy $U$ in the Hartree-Fock Hamiltonian is local, one immediately obtains the transport equation with collision term for proton ($q = p$) and neutron ($q = n$) distribution functions $f_q$

$$\frac{\partial}{\partial t} f_q - \{h_q + V_q, f_q\}_{\vec{r}, \vec{p}} = K(f_q) \hspace{1cm} (2)$$

or

$$\frac{\partial}{\partial t} f_q + \vec{v}_q \cdot \vec{\nabla}_r f_q - \vec{\nabla}_r (U_q + V_q) \cdot \vec{\nabla}_p f_q = K(f_q) \hspace{1cm} (3)$$

where $h = T + U$ is the mean-field Hamiltonian, $T$ is the kinetic energy, $V$ is the external field, and $K(f)$ is the non-Markovian collision term. When $f$ and $U$ change by a small amount around the equilibrium, we then have

$$f(\vec{r}, \vec{p}, t) = f_{eq}(\epsilon_p) + \delta f(\vec{r}, \vec{p}, t) \ , \quad U(\vec{r}, t) = U_0(\vec{r}) + \delta U(\vec{r}, t) \hspace{1cm} (4)$$

with

$$f_{eq}(\epsilon_p) = \frac{1}{(1 + e^{\beta(\epsilon_p - \mu)})} \ , \quad \delta U(\vec{r}, t) = \left( \frac{\partial U}{\partial \rho} \right)_{\rho_0} \delta \rho(\vec{r}, t) \ , \quad \epsilon_p = \frac{p^2}{2m} \ .$$  \hspace{1cm} (5)

The equation of motion of the small amplitude vibrations in the semi-classical limit for infinite nuclear matter is then obtained as
\[
\frac{\partial}{\partial t} \delta f + \vec{v} \cdot \nabla_r \delta f - \nabla_r [\delta U + 2\delta V] \cdot \vec{v} f_{eq}(\epsilon_p) = \delta K
\] (6)

which is the well-known the linearized Landau-Vlasov equation with a collision term. In this equation \( \delta f \equiv \delta f_n - \delta f_p \), \( \delta U \equiv \delta U_n - \delta U_p \) and \( \delta V \equiv \delta V_n - \delta V_p \) (\( \delta V_{n,p} = \pm \delta V \)) are differences between the indicated neutron and proton functions.

The isovector mean field \( \delta U(\vec{r}, t) \) can be expressed as \([7,11]\)

\[
\delta U(\vec{r}, t) = f_0 \delta \rho(\vec{r}, t)
\] (7)

where \( f_0 = F'_0(T)/N(T) \) is the quasiparticle zero-order interaction amplitude, \( F'_0(T) \) is the isovector Landau parameter

\[
F'_0(T) \simeq F'_0(T = 0) \left[ 1 - \frac{\pi^2}{12} \left( \frac{T}{\epsilon_F} \right)^2 \right],
\] (8)

\[
\delta \rho(\vec{r}, t) = \int \frac{g d\vec{p}}{(2\pi \hbar)^3} \delta f(\vec{r}, \vec{p}, t)
\] (9)

is the density distribution function, \( g = 2 \) is the spin degeneracy factor and

\[
N(T) = \int \frac{g d\vec{p}}{(2\pi \hbar)^3} \left( -\frac{\partial f_{eq}(\epsilon_p)}{\partial \epsilon_p} \right)
\] (10)

is the thermally averaged density of states. For \( T = 0 \) \( N(0) \) is given as \( N(0) = g p_F m/2\pi^2 \hbar^3 \), with \( p_F \) Fermi momentum.

Now, we present the details of the calculation of GDR response function for infinite and finite nuclear matter by using the linearized Landau-Vlasov equation without and with a non-Markovian collision term. The solution of Eq. (11) can be found in form of a plane wave for infinite nuclear matter

\[
\delta f(\vec{r}, \vec{p}, t) = \delta f_{k,\omega}(\vec{p}) e^{i(\vec{k} \cdot \vec{r} - (\omega + i\eta)t)},
\] (11)

where \( \eta \) is the vanishingly small positive number corresponding to an adiabatic switching of the field at time \( t = -\infty \). From the collisionless Landau-Vlasov equation, we have

\[
\frac{\delta f(\vec{r}, \vec{p}, t)}{\omega + i\eta - \vec{k} \cdot \vec{v}} + \frac{\vec{k} \cdot \vec{v}}{\omega + i\eta - \vec{k} \cdot \vec{v}} \frac{\partial f_{eq}(\epsilon_p)}{\partial \epsilon_p}[\delta U + 2\delta V] = 0 ,
\] (12)
and then integrating \( \int \frac{gdp}{(2\pi\hbar)^3} \) with weight 1, we obtain

\[
\delta \rho + [f_0 \delta \rho + 2\delta V] \chi^{(1)}(\vec{k}, \omega) = 0 ,
\]

where

\[
\chi^{(1)}(\vec{k}, \omega) = \int \frac{gdp}{(2\pi\hbar)^3} \frac{\vec{k}.\vec{v}}{\omega + i\eta - \vec{k}.\vec{v}} \frac{\partial f_{eq}(\epsilon_p)}{\partial \epsilon_p}
\]

is the unperturbed Lindhard function. We can then write response of the collisionless system for an external field \( \delta V \propto e^{i[\vec{k}.\vec{r}-(\omega+i\eta)t]} \) as

\[
\Pi^0(\vec{k}, \omega) = -\frac{\delta \rho}{\delta V} = \frac{2\chi^{(1)}(\vec{k}, \omega)}{1 + f_0 \chi^{(1)}(\vec{k}, \omega)} .
\]

Performing integrations in Eq.(14), we can find the real and imaginary parts of \( \chi^{(1)}(\vec{k}, \omega) \) as (for details please refer to [7,11])

\[
Im \chi^{(1)}(\vec{k}, \omega) = -\frac{\pi}{2} N(0)s \left[ \frac{m\omega}{kp_F} \right] f_{eq}(s^2 \bar{\epsilon})
\]

and

\[
Re \chi^{(1)}(\vec{k}, \omega) = \frac{N(0)}{4} \left[ \frac{m\omega}{kp_F} \right] f_{eq}(s^2 \bar{\epsilon}) .
\]

Here,

\[
\bar{\epsilon} = \frac{5}{3\rho_{eq}} \int \frac{gdp}{(2\pi\hbar)^3} \epsilon_p f_{eq}(\epsilon_p) ,
\]

\[
\rho_{eq} = \int \frac{gdp}{(2\pi\hbar)^3} f_{eq}(\epsilon_p)
\]

are quasiparticle average kinetic energy and density, respectively and

\[
s = \frac{m\omega}{kp_F} \left( \frac{e_F}{\bar{\epsilon}} \right)^{1/2} .
\]

The strength distribution function is obtained from the imaginary part of the response function [15]

\[
S(\vec{k}, w) = -\frac{1}{\pi} Im \Pi^0(\vec{k}, \omega) = -\frac{1}{\pi} \frac{2Im \chi^{(1)}(\vec{k}, \omega)}{(1 + f_0 Re \chi^{(1)}(\vec{k}, \omega))^2 + (f_0 Im \chi^{(1)}(\vec{k}, \omega))^2} .
\]
The solution of the linearized Landau-Vlasov equation with collision term for the infinite nuclear matter is given as

\[-i(\omega + i\eta)\delta f + i\vec{k}.\vec{v}\delta f - i\vec{k}.\vec{v}\frac{\partial f_{eq}(\epsilon_\rho)}{\partial \epsilon_\rho} [f_0\delta \rho + 2\delta V] = \delta K\, , \tag{22}\]

from which we obtain

\[\delta \rho + [f_0\delta \rho + 2\delta V]\chi^{(1)}(\vec{k}, \omega) = -[f_0\delta \rho + 2\delta V]\chi^{(2)}(\vec{k}, \omega)\, . \tag{23}\]

The collisional response function \(\chi^{(2)}(\vec{k}, \omega)\) can be expressed as [12]

\[\chi^{(2)}(\vec{k}, \omega) = \frac{1}{(2\pi)^3} \int d^3p_1 d^3p_2 d^3p_3 d^3p_4 \left(\frac{\Delta Q}{2}\right)^2 \frac{W(12;34)}{\pi} \frac{f_1 f_2 \overline{f}_3 \overline{f}_4 - f_1 \overline{f}_2 f_3 f_4}{w - \epsilon_3 - \epsilon_4 + \epsilon_1 + \epsilon_2 + i\eta} \tag{24}\]

where \(f_i = f_{eq}(\epsilon_i), \overline{f}_i = 1 - f_{eq}(\epsilon_i), \Delta Q = Q_1 + Q_2 - Q_3 - Q_4\) with \(Q_i = 1/|w - \vec{k}.\vec{v}_i|\), \(\epsilon_i = (m/2)v_i^2\) and \(W(12;34)\) denotes the basic two-body transition rate

\[W(12;34) = \frac{\pi}{(2\pi\hbar)^6} |v| < \frac{\vec{p}_1 - \vec{p}_2}{2} > |^2 \delta^3(\vec{p}_1 + \vec{p}_2 - \vec{p}_3 - \vec{p}_4)\, . \tag{25}\]

which can be expressed in terms of the scattering cross-section as

\[W(12;34) = \frac{1}{(2\pi\hbar)^3} \frac{4\hbar}{m^2} \frac{d\sigma}{d\Omega}_{pn} \delta^3(\vec{p}_1 + \vec{p}_2 - \vec{p}_3 - \vec{p}_4)\, . \tag{26}\]

Then the retarded response function with collision term is obtained as

\[\Pi^{coll}(\vec{k}, \omega) = -\frac{\delta \rho}{\delta V} = \frac{2\chi(\vec{k}, \omega)}{1 + f_0\chi(\vec{k}, \omega)} \tag{27}\]

with \(\chi(\vec{k}, \omega) = \chi^{(1)}(\vec{k}, \omega) + \chi^{(2)}(\vec{k}, \omega)\). Thus we can rewrite the strength distribution for our system with collision term as

\[S(\vec{k}, \omega) = -\frac{1}{\pi} Im\Pi^{coll}(\vec{k}, \omega) = -\frac{1}{\pi} \frac{2Im\chi(\vec{k}, \omega)}{(1 + f_0Re\chi(\vec{k}, \omega))^2 + (f_0Im\chi(\vec{k}, \omega))^2} \tag{28}\]

The strength function satisfies the following energy weighted sum rule [10]

\[\int_0^\infty d\omega \omega S(\vec{k}, \omega) = \frac{k^2}{2m\rho_0} \tag{29}\]

where \(\rho_0 = 0.16 fm^{-3}\) is the saturation density of infinite nuclear matter.
In the calculations of the collisional response function, we use conservation laws and symmetry properties. It is possible to reduce the twelve dimensional integrals to five fold integrals by the total and relative momentum transformations ($\vec{P} = \vec{p}_1 + \vec{p}_2$, $\vec{P}' = \vec{p}_3 + \vec{p}_4$, 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. We neglect the real part $\text{Re}\chi^{(2)}(\vec{k}, \omega)$ of the function $\chi^{(2)}(\vec{k}, \omega)$ in our calculations. The collisional response function $\chi^{(2)}(\vec{k}, \omega)$ has a singular behavior arising from the pole of the distortion functions, $Q_i = 1/[w - \vec{k} \cdot \vec{v}_i]$. We avoid this singular behavior by incorporating a pole approximation. In the distortion functions, we make the replacement $\omega \rightarrow \omega_D - i\Gamma/2$ where $\omega_D$ and $\Gamma$ are determined from $1 + f_0\chi^{(1)}(\vec{k}, \omega) = 0$ at each temperature that is considered. So, we evaluate the remaining five dimensional integrals numerically by employing a fast algorithm. In the evaluation of momentum integrals we make the replacement $(d\sigma/d\Omega)_{pn} \rightarrow \sigma_{pn}/4\pi$ with $\sigma_{pn} = 40$ mb, thus neglecting the angular anisotropy of the cross section.

In order to apply our results to nuclear dipole vibrations and finite nuclei, we work within the framework of Steinwedel and Jensen model which describes the GDR in heavy nuclei as a volume polarization mode conserving the total density $\rho_0 = \rho_n + \rho_p$ for infinite nuclear matter [16] where neutron and proton oscillate inside a sphere of radius $R$ as

$$\rho_p(\vec{r}, t) - \rho_n(\vec{r}, t) \propto \sin(\vec{k} \cdot \vec{r}) e^{i\omega t}. \quad (30)$$

According to this model, we choose the wave number of the normal mode as $k = \pi/2R$. 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 $R = 1.13A^{1/3}$.

The Landau parameter $F_0'(T = 0)$ can be expressed as a function of the symmetry energy coefficient $a_\tau$ in the Weizscker mass formula at zero temperature as follows [17]

$$F_0'(T = 0) = \frac{3a_\tau}{\epsilon_F} - 1. \quad (31)$$

For the value of $a_\tau = 28$ $MeV$ we have $F_0'(T = 0) = 1.33$. The value of the $F_0'(T)$ decreases with temperature because of the decrease of the thermally averaged level density $N(T)$ [7].

So far, our GDR calculations have been for infinite nuclear matter. In the rest of the paper we will employ the Thomas-Fermi approximation (TF) to calculate GDR response
function by using the linearized Landau-Vlasov equation with and without the collision term at finite temperature for finite nuclei. The Thomas-Fermi theory, together with its extensions, is the semiclassical treatment of nuclear dynamics in its independent particle or Hartree-Fock approximation and can be explained from quite different points of view [15,16]. We evaluate the expression for the GDR in the TF approximation, which corresponds to a semi-classical transport description of the collective vibrations. The solution of the linearized Landau-Vlasov equation with and without collision term

$$\frac{\partial}{\partial t} \delta f + \vec{v} \cdot \nabla_r \delta f - \nabla_r [\delta U + 2\delta V] \cdot \nabla_p f_{eq}(\epsilon_p, r) = \delta K \tag{32}$$

for finite nuclei is obtained with the local plane wave ansatz

$$\delta f(r, p, t) = \delta f_{k,\omega}(r, p) e^{i(k \cdot r - (\omega + i\eta)t)} \tag{33}$$

Response of the finite system without and with collision term is then obtained as

$$\Pi^0_{TF}(\vec{k}, \omega) = \frac{2\chi^{(1)}_{TF}(\vec{k}, \omega)}{1 + f^0_{TF}\chi^{(1)}_{TF}(\vec{k}, \omega)} , \quad \Pi^{\text{coll}}_{TF}(\vec{k}, \omega) = \frac{2\chi_{TF}(\vec{k}, \omega)}{1 + f^0_{TF}\chi_{TF}(\vec{k}, \omega)} \tag{34}$$

where $\chi_{TF}(\vec{k}, \omega) = \chi^{(1)}_{TF}(\vec{k}, \omega) + \chi^{(2)}_{TF}(\vec{k}, \omega)$. The strength distribution function without and with collision term are

$$S_{TF}(\vec{k}, w) = -\frac{1}{\pi} Im \Pi^0_{TF}(\vec{k}, \omega) , \quad S_{TF}(\vec{k}, w) = -\frac{1}{\pi} Im \Pi^{\text{coll}}_{TF}(\vec{k}, \omega) \tag{35}$$

with

$$\chi^{(i)}_{TF}(\vec{k}, \omega) = \frac{1}{A} \int d\vec{r} \rho(r) \chi^{(i)}(\vec{k}, \omega, r) \tag{36}$$

where $i = 1, 2$ , and

$$f^0_{TF} = \frac{1}{A} \int d\vec{r} \rho(r) f_0(r) \tag{37}$$

The function $\chi^{(2)}(\vec{k}, \omega, r)$ is obtained by evaluating the collision term given in Eq. (24) using Thomas-Fermi approximation. We determine the nuclear density $\rho(r)$ for the finite nuclear matter in TF approximation using a Wood-Saxon potential with a depth $V_0 = 44 \text{ MeV}$, thickness parameter $t_p = 0.67 \text{ fm}$ and sharp radius $R = 1.13A^{1/3} [16]$,
\[ V(r) = -\frac{V_0}{(1 + e^{r/R})}, \quad \rho(r) = \frac{2}{3\pi^2} k_F(r)^3 \Theta(\lambda - V(r)), \quad k_F(r) = \left(\frac{2m}{\hbar^2}[V(r_c) - V(r)]\right)^{1/2} \]

where \( r_c \) is the critical radius for a mass number \( A \) is defined as

\[ A = \int_0^{r_c} \rho(r) \, dr \]  \hspace{1cm} (39)

Here \( \lambda = V(r_c) \) and, the expression for \( A \) can be numerically integrated for a given \( A \) to determine \( r_c \). For finite nuclear matter, the interaction amplitude is \( f_0(r) = 3V_0'(r)\rho(r)/(2\epsilon_F(r) N(r, T = 0)) \) which is related to the parameters of the simplified Skyrme force as

\[ V_0'(r) = -\frac{1}{2} t_0 (x_0 + \frac{1}{2}) - \frac{1}{8} t_3 \rho(r) \]  \hspace{1cm} (40)

We use the following parameters: \( t_0 = -983.4 \text{ MeV} \text{fm}^3, t_3 = 13106 \text{ MeV} \text{fm}^6, \) and \( x_0 = 0.48 \) [10].

We show our results for the GDR strength function with and without the collision term in Fig. 1 and in Fig. 2 for \(^{120}\text{Sn}\) and \(^{208}\text{Pb}\), respectively, calculated using infinite nuclear matter formalism within the framework of Steinwedel-Jensen model. In Fig. 3 and in Fig. 4 we show the GDR strength function with and without collision term for \(^{120}\text{Sn}\) and \(^{208}\text{Pb}\), respectively, calculated employing Thomas-Fermi approximation. In this figures, we also compare our results with the normalized experimental data taken from [1]. The temperature parameter \( T \) in the mean occupation number functions \( f(\epsilon, T) \) is related to the experimental temperature \( T^* \) as \( T = T^* \sqrt{a_E/a_F} \) where \( a_F = A\pi^2/4\epsilon_F \) is the Fermi gas level density parameter and \( a_E \) is the energy dependent empirical level density parameter [1].

From these figures, we first note that without the collision term the position of the peak of the strength function does not change appreciably with temperature. This behavior is in accordance with the experimental results [1,2]. Since we neglect the real part of the collisional response, when we include the collisional term the average position of the peak values of the strength functions do not change again, but as the result of the collisions the
overall shape of the strength function changes somewhat and this change becomes more pronounced with increasing temperature. In the case of infinite nuclear matter this change has the tendency to improve the agreement with experimental results for $^{120}Sn$ but this tendency is much less pronounced for $^{208}Pb$ as it can be seen by comparing Fig. 1 with Fig. 2. On the other hand, for this case of finite nuclei calculations employing Thomas-Fermi approximation the change produced by the addition of the collision term can be clearly noted in Fig. 3 for $^{120}Sn$ and in Fig. 4 for $^{208}Pb$. Indeed, in both cases, the overall agreement with the experimental results is much more improved when the binary collision term is included.

In our work, we obtain a reasonable description of the giant dipole excitations in $^{120}Sn$ and $^{208}Pb$ using semi-classical approach with Thomas-Fermi approximation as compared to infinite nuclear matter formalism, and we demonstrate the importance of the collision term which improves the agreement of the calculated strength functions with the experimental results. We believe that inclusion of the coherent damping mechanism into our formalism will extend our description further.
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FIG. 1. The GDR strength function of $^{120}\text{Sn}$ obtained using Steinwedel-Jensen model. Solid and dashed lines show the response function without and with the collision term, respectively. The normalized data is taken from [1].
FIG. 2. The GDR strength function of $^{208}$Pb obtained using Steinwedel-Jensen model. Solid and dashed lines show the response function without and with the collision term, respectively. The normalized data is taken from [1].
FIG. 3. The GDR strength function of $^{120}$Sn calculated by Thomas-Fermi approximation. Solid and dashed lines show the response function without and with the collision term for the finite nuclear matter. The normalized data is taken from [1].
FIG. 4. The GDR strength function of $^{208}$Pb calculated by Thomas-Fermi approximation. Solid and dashed lines show the response function without and with the collision term for the finite nuclear matter. The normalized data is taken from [1].