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COLLISIONAL DAMPING OF GIANT RESONANCES WITH AN OPTIMIZED FINITE-RANGE EFFECTIVE INTERACTION

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Collisional damping widths of isoscalar giant monopole, isovector giant dipole and isoscalar giant quadrupole excitations for ^{120}Sn and ^{208}Pb at zero and finite temperatures are calculated within Thomas-Fermi approximation by employing a recently proposed modified Skyrme force and the results are compared with previous works.

Keywords: Giant resonances; collisional damping; effective interaction.

1. Introduction

The study of collisional damping of giant resonances at zero and finite temperatures is an active area of investigation.[1–4] Theoretical studies of nuclear collective vibrations built on the ground and the excited states of nuclei are mostly based on the Random Phase Approximation (RPA) theory. However, since RPA approach is the small amplitude limit of the Time Dependent Hartree Fock Theory (TDHF), it is not suitable for describing the damping of the collective excitations.[5] In order to describe nuclear collective response including damping, it is necessary to incorporate coupling between the collective states and the doorway configurations.[6]

Recently, the collisional damping widths of giant monopole vibrations, giant dipole vibrations and giant quadrupole vibrations at zero and finite temperature were calculated for different mass numbers in Thomas–Fermi approximation by employing the microscopic in medium cross-sections of Li and Machleidt and by using the phenomenological zero range Skyrme force and the finite range Gogny force.[7,8] In these works it was reported that the cross-sections based on the phenomenological zero range Skyrme force and finite range Gogny force exhibit unrealistic behaviour as a function of density which results in larger damping. In order to incorporate the finite range effects into effective interactions, a modification of the Skyrme force by introducing a cut-off factor for high momentum transfer is recently proposed.[9] In this work, we use this new proposed effective Skyrme force with cut off to calculate the collisional damping of giant resonances for ^{120}Sn and for ^{208}Pb at zero and at finite temperatures.

2. Calculations and Results

In Ref. 3 the collisional relaxation rates of giant vibrations in nuclei are calculated using a semiclassical transport equation with a non-Markovian collision term. The isovector and the isoscalar decay widths are given as $\Gamma_\lambda = \int d\mathbf{r}\Gamma_\lambda(r)$, with

$$\Gamma_\lambda^v = \frac{1}{N_\lambda}\int d\mathbf{p}_1 d\mathbf{p}_2 d\mathbf{p}_3 d\mathbf{p}_4 \left[(W_{pp}+W_{nn})\left(\frac{\Delta\chi_\lambda}{2}\right)^2 + 2W_{pn}\left(\frac{\widetilde{\Delta\chi_\lambda}}{2}\right)^2\right] Z f_1 f_2 \bar{f}_3 \bar{f}_4 \tag{1}$$

and a similar expression for the isoscalar modes,

$$\Gamma_\lambda^s = \frac{1}{N_\lambda}\int d\mathbf{r} d\mathbf{p}_1 d\mathbf{p}_2 d\mathbf{p}_3 d\mathbf{p}_4 [W_{pp}+W_{nn}+2W_{pn}]\frac{(\Delta\chi_\lambda)^2}{2} Z f_1 f_2 \bar{f}_3 \bar{f}_4\,, \tag{2}$$

where $N_\lambda = \int d\mathbf{r} d\mathbf{p}(\chi_\lambda)^2(-\partial f/\partial\epsilon)$ is a normalization, f_i corresponds to the Fermi–Dirac distribution function and

$$\bar{f}_i = 1 - f_i, \Delta\chi_\lambda = \chi_\lambda(1) + \chi_\lambda(2) - \chi_\lambda(3) - \chi_\lambda(4)\,,$$

$$\widetilde{\Delta\chi_\lambda} = \chi_\lambda(1) - \chi_\lambda(2) - \chi_\lambda(3) + \chi_\lambda(4)\,,$$

$$Z = [\delta(\hbar\omega_\lambda - \Delta\epsilon) - \delta(\hbar\omega_\lambda + \Delta\epsilon)]/\hbar\omega_\lambda\,,$$

and $\chi_\lambda(t)$ denotes the distortion factor of the phase-space density $\delta f(t) = \chi_\lambda(t)(-\partial f/\partial\epsilon)$ in the corresponding mode. In this expression, transition rates W_{pp}, W_{nn}, W_{pn} associated with proton-proton, neutron-neutron and proton-neutron collisions are given in terms of the corresponding scattering cross-sections as

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

For the in medium nucleon nucleon cross-section we use the expression of Ref. 9, which is given as

$$\left(\frac{d\sigma}{d\Omega}\right)_{xy} = \frac{1}{4\pi}\mathcal{R}_{xy}(\rho)\left[\overline{\mathcal{C}}(|\langle\mathbf{q}^2\rangle|)\right]^2, \tag{4}$$

where $\mathcal{R}_{xy}(\rho)$ includes Skyrme force parameters and the labels xy denote the cross-section in the proton-neutron channel (np), the proton-proton channel (pp) and the neutron-neutron channel (nn), which are given by

$$\mathcal{R}_{pn}(\rho) = \frac{\pi m^2}{4\hbar(2\pi\hbar)^3}\frac{1}{2}\left(\left[t_0(1-x_0)+\frac{t_3}{6}(1-x_3)\rho^\alpha\right]^2 + 3\left[t_0(1+x_0)+\frac{t_3}{6}(1+x_3)\rho^\alpha\right]^2\right) \tag{5}$$

and

$$\mathcal{R}_{pp}(\rho) = \frac{\pi m^2}{4\hbar(2\pi\hbar)^3}\left[t_0(1-x_0) + \frac{t_3}{6}(1-x_3)\rho^\alpha\right]^2 \tag{6}$$

with a similar expression for the neutron-neutron cross-section. Here $\overline{\mathcal{C}}$ is the cut-off factor:

$$\overline{\mathcal{C}}(q,q') = \frac{1+\beta_1^2\frac{q^2+q'^2}{\hbar^2}}{1+\beta_2^2\frac{q^2+q'^2}{\hbar^2}}\,, \tag{7}$$

where $\mathbf{q} = (\mathbf{p}_1 - \mathbf{p}_2)/2$ and $\mathbf{q}' = (\mathbf{p}_3 - \mathbf{p}_4)/2$ represent the relative momenta before and after collisions.

We now present the calculations of the collisional damping widths of isovector giant dipole resonance (GDR), isoscalar giant monopole resonance (GMR) and isoscalar giant quadrupole resonance (GQR) in some details. Although the physical picture and the mathematical formulation of each giant resonance is different, our calculation scheme is a generalized one, which can be applied to all three resonances. The beginning point of our calculation scheme is a coordinate transformation that allows us to evaluate the momentum integrals exactly which were treated only approximately previously.[3] We apply formula (1) to calculate the collisional damping width of the GDR, where due to momentum conservation, terms involving W_{pp} and W_{nn} drop out, and the damping is determined by the proton-neutron collision term. Making a change of variables to the relative momentum $\mathbf{q} = \frac{1}{2}(\mathbf{p_1} - \mathbf{p_2}), \mathbf{q}' = \frac{1}{2}(\mathbf{p_1} - \mathbf{p_2})$ and total momentum $\mathbf{K} = (\mathbf{p_1} + \mathbf{p_2})$, $\mathbf{K}' = (\mathbf{p_3} + \mathbf{p_4})$ before and after the collision respectively, the expression for the damping width in terms of the new variables becomes

$$\Gamma_D^v = \frac{1}{N_D}\int d\mathbf{r}d\mathbf{K}d\mathbf{K}'d\mathbf{q}d\mathbf{q}' 2W_{pn}\left(\frac{\widetilde{\Delta\chi_D}}{2}\right)^2 Zf_1f_2\bar{f}_3\bar{f}_4\,. \tag{8}$$

We take the distortion factors according to the scaling picture as $\chi_D(\mathbf{p}) = \mathbf{p}\cdot\mathbf{e}$, where $\mathbf{e}$ denotes the unit vector along the relative displacement of the proton and neutron momentum distributions which is equivalent to the Legendre function parametrization, $\chi_D = pP_1(\cos\theta)$ of Ref. 3. The distortion function is

$$\left(\frac{\widetilde{\Delta\chi_D}}{2}\right)^2 = (\mathbf{q}\cdot\mathbf{e} - \mathbf{q}'\cdot\mathbf{e})^2\,, \tag{9}$$

which can be effectively written as

$$(\mathbf{q}\cdot\mathbf{e} - \mathbf{q}'\cdot\mathbf{e})^2 \to \frac{1}{3}(\mathbf{q}^2 + \mathbf{q}'^2)\,. \tag{10}$$

This is due to the fact that the integral of the cross terms vanishes and the integral is independent of the orientation of the unit vector $\mathbf{e}$. We also make another replacement:

$$Z = \frac{1}{\hbar\omega_D}[\delta(\hbar\omega_D - \Delta\epsilon) - \delta(\hbar\omega_D + \Delta\epsilon)] \to \frac{1}{\hbar\omega_D}(1 - e^{-\frac{\hbar\omega_D}{T}})\delta(\hbar\omega_D - \Delta\epsilon)\,. \tag{11}$$

Putting these altogether and carrying out the integral over $d\mathbf{K}'$ with the momentum conserving Dirac delta function, we have the expression for the damping width of the dipole resonance in a closed form as

$$\Gamma_D^v = \frac{1}{N_D}\int d\mathbf{r}d\mathbf{K}d\mathbf{q}d\mathbf{q}'\left(2\frac{1}{(2\pi\hbar)^3}\frac{4\hbar}{m^2}\left(\frac{d\sigma}{d\Omega}\right)_{pn}\frac{1}{3}(\mathbf{q}^2+\mathbf{q}'^2)\right.$$
$$\left.\times\frac{1}{\hbar\omega_D}(1-e^{-\frac{\hbar\omega_D}{T}})\delta(\hbar\omega_D-\Delta\epsilon)f_1f_2\bar{f}_3\bar{f}_4\right). \tag{12}$$

In the further evaluation of the momentum integrals, we neglect the angular anisotropy of the cross-sections and make the replacement $(d\sigma/d\Omega)_{pn}\rightarrow\sigma_{pn}/4\pi$ and we obtain the damping width as

$$\Gamma_D = \frac{2}{N_D}\int d^3r\left[\frac{1}{(2\pi\hbar)^3}\frac{4\hbar}{m^2}\right]\frac{1}{\hbar\omega_D}\left[1-\exp\left(-\frac{\hbar\omega_D}{T}\right)\right]I_D(r)\,, \tag{13}$$

with

$$I_D(r) = \int d^3Kd^3qd^3q'\frac{1}{3}(\mathbf{q}^2+\mathbf{q}'^2)\left(\frac{\sigma_{pn}}{4\pi}\right)\delta(\hbar\omega_D-\epsilon'+\epsilon)f_1f_2\bar{f}_3\bar{f}_4\,, \tag{14}$$

where $\epsilon=\mathbf{q}^2/m$ and $\epsilon'=\mathbf{q}'^2/m$ represent energies of two particle system in the center of mass frame before and after the binary collision. This expression then becomes

$$I_D(r) = (2\pi)^2(m\sqrt{m})^3\frac{m}{3}\left[\int_0^\infty\int_0^\infty\sqrt{E}dE\sqrt{\epsilon\epsilon'}d\epsilon\sigma_{pn}(\epsilon+\epsilon')\right.$$
$$\left.\times\int_{-1}^{+1}\int_{-1}^{+1}dzdz'f_1f_2\bar{f}_3\bar{f}_4\right], \tag{15}$$

where $\epsilon'=\epsilon+\hbar\omega_D$ and $E=K^2/4m$ denotes the kinetic energy of the center of mass. In terms of the integration variables, the products of the phase-space factors have the forms:

$$f_1f_2 = \frac{1}{1+A^2(E,\epsilon)+A(E,\epsilon)\Big(\exp\left(\frac{+z\sqrt{E\epsilon}}{T}\right)+\exp\left(\frac{-z\sqrt{E\epsilon}}{T}\right)\Big)}\,, \tag{16}$$

and

$$\bar{f}_3\bar{f}_4 = \frac{A^2(E,\epsilon')}{1+A^2(E,\epsilon')+A(E,\epsilon')\Big(\exp\left(\frac{+z'\sqrt{E\epsilon'}}{T}\right)+\exp\left(\frac{-z'\sqrt{E\epsilon'}}{T}\right)\Big)}\,, \tag{17}$$

where $A(E,\epsilon)=\exp((E+\epsilon)/2-\epsilon_F)/T$, and $z=\cos\theta$ and $z'=\cos\theta'$ are the angles between $\mathbf{q}$ and $\mathbf{K}$, and $\mathbf{q}'$ and $\mathbf{K}$, respectively. It is possible to carry out the angular integrals over z and z' analytically:

$$\int_{-1}^{+1}\int_{-1}^{+1}dzdz'f_1f_2\bar{f}_3\bar{f}_4 = \left(\frac{1}{1+A(E,\epsilon)^2}X(A(E,\epsilon),\beta)\right)$$
$$\times\left(\frac{A(E,\epsilon+\hbar\omega_D)^2}{1+A(E,\epsilon+\hbar\omega_D)^2}\times X(A(E,\epsilon+\hbar\omega_D),\beta)\right),$$

where

$$\beta = \frac{\sqrt{E\epsilon}}{T} \tag{18}$$

and

$$X(A(E,\epsilon),\beta) = \frac{1}{\beta}\frac{A(E,\epsilon)^2+1}{A(E,\epsilon)^2-1}\,\mathrm{Log}\left[e^{2\beta}\left(\frac{A(E,\epsilon)+e^{-\beta}}{A(E,\epsilon)+e^{\beta}}\right)^2\right]. \tag{19}$$

The only expression left is the normalization which can be expresed as

$$N_D = \frac{(4\pi)^2}{3}\frac{4m^{5/2}}{\sqrt{2}}\int r^2 dr \int d\epsilon\epsilon^{3/2}\frac{e^{(\epsilon-\mu)/T}}{[e^{(\epsilon-\mu)/T}+1]^2 T}. \tag{20}$$

Figure 1 shows the collisional damping width of GDR in ^{120}Sn and ^{208}Pb as a function of the experimental temperature and comparison with data. Calculations performed with the Skyrme force with cut-off $(\mathrm{SkM}^*)^{\mathrm{cut}}$ cross-sections are indicated with long-dashed lines. For comparison, we also include the results of Ref. 6 with the SkM* (solid lines) and the Gogny (dashed lines) cross-sections with the bare nucleon mass, and the points are the experimental data from Ref. 2. The results with the $(\mathrm{SkM}^*)^{\mathrm{cut}}$ cross-section gives about 30%–35% of the experimental damping widths in tin and lead nuclei at finite temperatures.

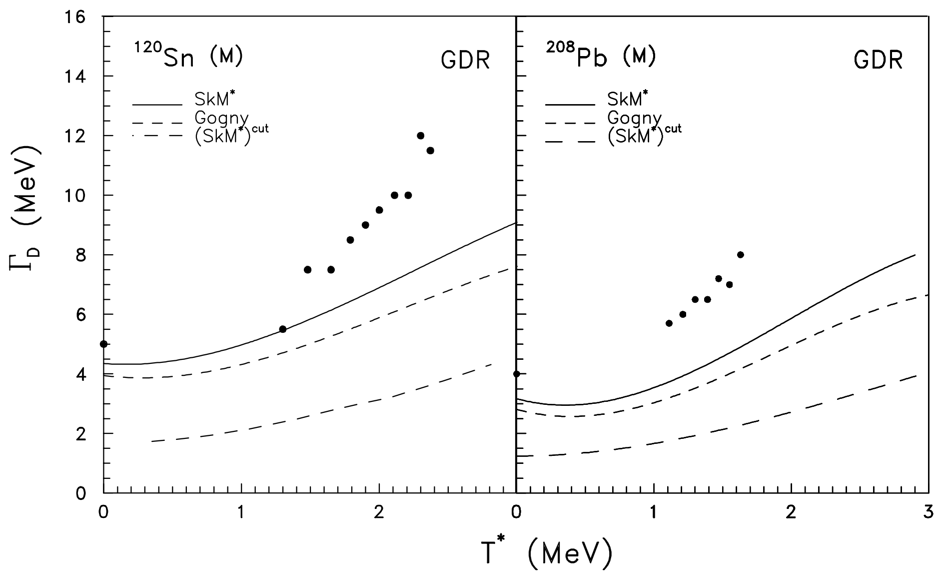

Fig. 1. The collisional damping width of GDR in ^{120}Sn and ^{208}Pb as a function of temperature. Long dashed lines are results with the cross-sections of $(\mathrm{SkM}^*)^{\mathrm{cut}}$, and solid lines are with the SkM* and dashed lines are with the Gogny cross-sections with the bare nucleon mass. The data is taken from Ref. 2.

In order to calculate the collisional damping width of the isoscalar GMR and isoscalar GQR we apply the formula (2), where we use the nuclear hydrodynamics model to calculate the distortion factors for GMR and GQR as well as the scaling model for GQR in order to make comparison with previous works. In the nuclear hydrodynamics model the distortion factors of the momentum distribution are expressed in terms of the velocity field $\Phi(r)$ associated with the collective mode as $\chi = (\mathbf{p}\cdot\nabla)(\mathbf{p}\cdot\nabla)\Phi(r)$.[10] An accurate description of the monopole and quadrupol vibrations can be obtained by parametrizing the velocity field in terms of the zeroth order Bessel functions $\Phi(r) = j_0(kr)$, with the wave number $k = \pi/R$, and in terms of the second order Bessel functions $\Phi(r) = j_2(kr)$, with the wave number $k = 3.3421/R$,[11] respectively where R is the the nuclear radius. The distortion function for each momentum is

$$\chi_\alpha = \left[j''_\alpha(kr) - \frac{1}{kr}j'_\alpha(kr)\right](\mathbf{p}\cdot\mathbf{r})^2 + \frac{1}{kr}j'_\alpha(kr)p^2\,, \tag{21}$$

where the primes indicate derivatives with respect to arguments and the label α stands for the monopole/quadrupole $\alpha = 0/\alpha = 2$ cases, and the distortion factor is expressed as

$$\Delta\chi_\alpha = \chi_\alpha^{(1)} + \chi_\alpha^{(2)} - \chi_\alpha^{(3)} - \chi_\alpha^{(4)}\,. \tag{22}$$

In the case of isoscalar modes, the collisional width is determined by the spin-isospin averaged nucleon nucleon cross-section, $(d\sigma/d\Omega)_0 = [(d\sigma/d\Omega)_{pp} + (d\sigma/d\Omega)_{nn} + 2(d\sigma/d\Omega)_{pn}]/4$,[3] and making the replacement $(d\sigma/d\Omega)_0 \to \sigma_0/4\pi$ we obtain the damping width for GMR and GQR as

$$\begin{aligned}
\Gamma_\alpha(r) = {} & \left(\frac{1}{(2\pi\hbar)^3}\frac{4\hbar}{m^2}\frac{(1-e^{-\frac{\hbar\omega_\alpha}{T}})}{\hbar\omega_\alpha}(4\pi)(m\sqrt{m})^3\frac{1}{4}(2\pi)^2\right) \\
& \times \int dE d\epsilon \sqrt{E}\sqrt{\epsilon}\sqrt{\epsilon+\hbar\omega_\alpha}\left[\left(4m^2F^2(k,r)\epsilon^2\frac{4\pi}{5}\right.\right. \\
& + 4m^2F^2(k,r)(\epsilon+\hbar\omega_\alpha)^2\frac{4\pi}{5} + 4\pi 4m^2G^2(k,r) \\
& \left. + \; 8m^2F(k,r)G(k,r)\frac{4\pi}{3}\hbar\omega_\alpha\right)\left[\frac{1}{1+A(E,\epsilon)^2}X(A(E,\epsilon),\beta)\right] \\
& \times \left[\frac{A(E,\epsilon+\hbar\omega_\alpha)^2}{1+A(E,\epsilon+\hbar\omega_\alpha)^2}X(A(E,\epsilon+\hbar\omega_\alpha),\beta)\right] \\
& \left. + f(E,\epsilon,k,r)\right]\left(\frac{\sigma_0(E,\epsilon,r)}{4\pi}\right)\,,
\end{aligned} \tag{23}$$

where the function $f(E, \epsilon, k, r)$ is

$$f(E,\epsilon,k,r) = -4m^2F^2(k,r)2\epsilon(\epsilon+\hbar\omega_\alpha)\frac{4\pi}{15}\int_{-1}^{1}\int_{-1}^{1} dzdz'[3z^2z'^2 + z^2(1-z'^2) + (1-z^2)z'^2 + 2(1-z^2)(1-z'^2)]f_1f_2\bar{f}_3\bar{f}_4\,, \tag{24}$$

and we define $F(k,r)$ and $G(k,r)$ as

$$F(k,r) = \left(j''_\alpha(kr) - \frac{1}{kr}j'_\alpha(kr)\right)$$

$$G(k,r) = \left(\frac{1}{kr}j'_\alpha(kr)\hbar\omega_\alpha\right)\,. \tag{25}$$

The collisional damping width for both isoscalar GMR and GQR cases can be written as

$$\Gamma^s_\alpha = \frac{1}{N_\alpha}\int r^2\Gamma_\alpha(r)dr\,, \tag{26}$$

where the normalization is

$$N_\alpha = 8(4\pi)(2\pi)m^{\frac{7}{2}}\sqrt{2}\int r^2dr\left\{\frac{1}{5}\left[j''_\alpha(kr) - \frac{1}{kr}j'_\alpha(kr)\right]^2 + \left(\frac{1}{kr}j'_\alpha(kr)\right)^2 + \frac{2}{3}\left[j''_\alpha(kr) - \frac{1}{kr}j'_\alpha(kr)\right]\left(\frac{1}{kr}j'_\alpha(kr)\right)\right\}\int \epsilon^{\frac{5}{2}}\frac{e^{(\epsilon-\mu)/T}}{T[e^{(\epsilon-\mu)/T}]^2}d\epsilon\,. \tag{27}$$

In the scaling model for the GQR the distortion factors of the momentum distribution are parametrized in terms of the Legendre function as[3]

$$\chi^S_Q = p^2P_2(\cos\theta)\,. \tag{28}$$

In a similar manner we obtain the collisional damping width expression using the scaling model distortion functions for GQR as

$$\Gamma^S_Q(r) = \frac{1}{(2\pi\hbar)^3}\frac{4\hbar}{m^2}\frac{(1-e^{-\frac{\hbar\omega_Q}{T}})}{\hbar\omega_Q}(4\pi)(m\sqrt{m})^3\frac{1}{4}(2\pi)^2 \times\int dEd\epsilon\sqrt{E}\sqrt{\epsilon}\sqrt{\epsilon+\hbar\omega_Q}f(E,\epsilon)\left(\frac{\sigma_0(E,\epsilon,r)}{4\pi}\right)\,, \tag{29}$$

where the function $f(E,\epsilon)$ is

$$f(E,\epsilon) = \left(\frac{m}{2}\right)^2\int_{-1}^{1}\int_{-1}^{1}dzdz'[6\epsilon z^2 + 6(\epsilon+\hbar\omega_Q)z'^2 + 2\hbar\omega_Q]^2f_1f_2\bar{f}_3\bar{f}_4\,. \tag{30}$$

The normalization is now expressed as

$$N^S_Q = \frac{4}{5}(4\pi)^2m^{\frac{7}{2}}\sqrt{2}\int r^2dr\int\epsilon^{\frac{5}{2}}\frac{e^{(\epsilon-\mu)/T}}{T[e^{(\epsilon-\mu)/T}]^2}d\epsilon\,. \tag{31}$$

The results for the calculations of collisional damping width of GMR and GQR in ^{120}Sn and ^{208}Pb as a function of temperature are shown in Figs. 2 and 3 where the upper panel shows the results in the hydrodynamic model and the lower panel shows the result in the scaling model respectively, and they are compared with experimental data at zero temperature.[9] Calculations performed with the $(\mathrm{SkM}^*)^{\mathrm{cut}}$ cross-sections are indicated with long-dashed lines. For comparison, we also include the results of Ref. 7 with the SkM* (solid lines) and the Gogny (dashed lines) cross-sections with the bare nucleon mass. The results with the $(\mathrm{SkM}^*)^{\mathrm{cut}}$ cross-section gives about 20%–30% of the experimental damping widths in tin and lead nuclei at zero temperature.

In the numerical evaluations, we determine the nuclear density $\rho(r)$ in Thomas–Fermi approximation using a Wood-Saxon potential with a depth $V_0 = -44$ MeV, thickness parameter $t_p = 0.67$ fm and sharp radius $R_0 = 1.27A^{1/3}$ fm, and in all calculations we use the position dependent chemical potential $\mu(r,T)$ in the Fermi–Dirac function $f(\epsilon, r, T) = 1/1 + e^{(\epsilon-\mu(r,T))/T}$ at each temperature. The chemical potential at each space point is determined in terms of the Fermi energy, and hence is found the nuclear density, at small but finite temperatures by inverting the equation of state of a noninteracting Fermi gas.[11] In calculating $(\mathrm{SkM}^*)^{\mathrm{cut}}$ cross-sections we use the SkM* parameters[7] $\alpha = 1/6$, $x_0 = 0.09$, $t_0 = -2645$ MeV fm^3, $t_3 = 15595$ MeV fm$^{7/2}$, $\beta_1 = \sqrt{0.29}$ fm and the cut-off paremeters $\beta_2 = \sqrt{1.66}$ fm.[8] For the mass dependence of the resonance energies we use the formulas

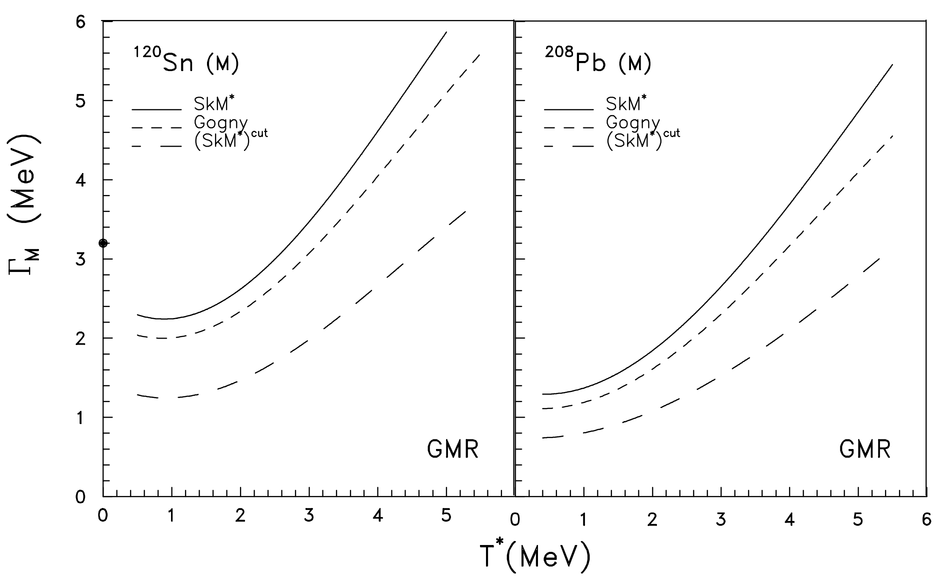

Fig. 2. The collisional damping width of GMR in ^{120}Sn and ^{208}Pb as a function of temperature. Long dashed lines are results with the cross-sections of $(\mathrm{SkM}^*)^{\mathrm{cut}}$, and solid lines are calculations with the SkM*, and dashed lines are results with the Gogny cross-sections with the bare nucleon mass.

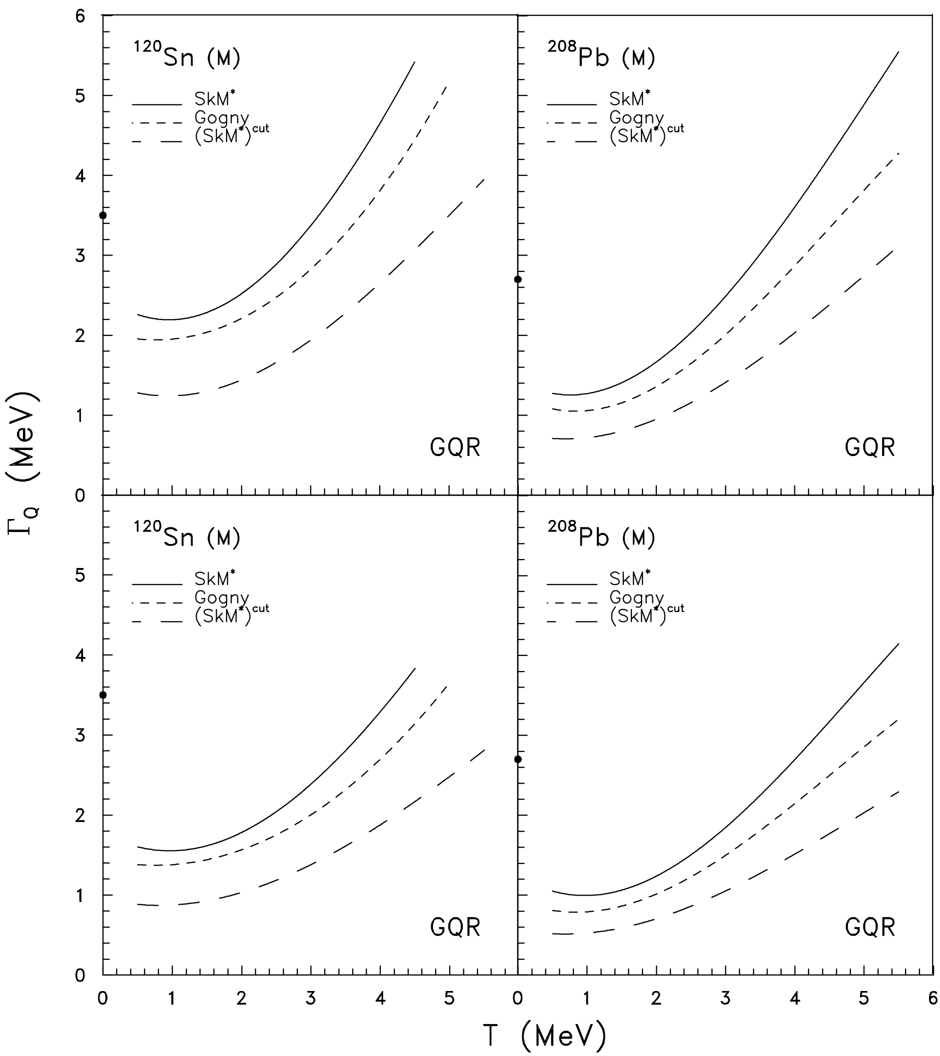

Fig. 3. The collisional damping width of GQR in ^{120}Sn and ^{208}Pb as a function of temperature. Long dashed lines are results with the cross-sections of (SkM*)$^{\rm cut}$, and solid lines are calculations with the SkM*, and dashed lines are results with the Gogny cross-sections with the bare nucleon mass. In the upper/lower panel for distortion factors hydrodynamic/scaling approximation is used.

$\hbar\omega_D = 80A^{-1/3}$ MeV for the GDR energies,[6] $\hbar\omega_M = 31.2A^{-1/3} + 20.6A^{-1/6}$ MeV to calculate the giant monopole resonance energies[7] and $\hbar\omega_Q = 64A^{-1/3}$ MeV for the giant quadrupole resonance energies.[7] In the Fermi gas picture, the effective excitation energy of the system and temperature are related according to

$$E^* = \int \frac{4}{(2\pi\hbar)^3} d^3r d^3p \epsilon [f(\epsilon, T) - f(\epsilon, T = 0)] \,. \tag{32}$$

For temperatures which are small compared to the Fermi energy $\epsilon_F = \mu(T = 0)$ and in sharp radius approximation, the relation between the excitation energy and temperature becomes, $E^* = a_F T^2$, where $a_F = A\pi^2/4\epsilon_F$ is the Fermi gas level density parameter. The experimental temperature T^* is determined from the excitation energy using a similar relation $T^* = \sqrt{E^*/a_E}$, where a_E denotes the energy dependent empirical level density parameter.[2] Hence, the temperature parameter in $f(\epsilon, T)$ is related to the experimental temperature as $T = T^*\sqrt{a_E/a_F}$.

3. Conclusion

In order to assess the fraction of the total width of collective excitations that is exhausted by decay into the incoherent 2p–2h states, namely the collisional damping, we need realistic in-medium nucleon-nucleon cross-sections around Fermi energy. In this respect this new proposed, modified Skyrme cross-section is a candidate for providing the best available input, since the cross-sections interpolate correctly between the free space and the medium at high momentum transfer. We can conclude that the collisional damping of the GDR, GMR and GQR excitations is not very strong, and accounts for about 1/3 of experimental results in both nuclei at zero and finite temperatures.

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