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Quantal and Thermal Dampings of the Hot Giant Dipole Resonance due to Complex Configuration Mixing

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An approach is presented to study the width of the giant dipole resonance (GDR) at non-zero temperature T, which includes all forward-going processes up to two-phonon ones. Calculations are performed in ^{120}Sn and ^{208}Pb. An overall agreement between theory and experiment is found. The total width of the GDR due to coupling of the GDR phonon to all ph, pp and hh configurations increases sharply as T increases up to $T \sim 3$ MeV and saturates at $T \sim 4$ - 6 MeV. The quantal width Γ_Q due to coupling to ph configurations decreases with increasing T. It is almost independent of T if the contribution of two-phonon processes at $T \neq 0$ is omitted.

The giant dipole resonance (GDR) built on compound nuclear states (hot GDR), has been one of major the subjects of this conference. (See also [1] for the reviews). The width of the GDR built on the ground state (g.s. GDR) is composed mainly of Landau splitting, the escape width $\Gamma^{\uparrow}$, and the spreading width $\Gamma^{\downarrow}$. The Landau splitting represents the distribution of the GDR over a number of harmonic oscillators with frequencies around the GDR energy. It amounts to only a small fraction of the total width. The escape width $\Gamma^{\uparrow}$ arises from a γ or particle emission. It can be included via coupling to the continuum and is also small (few hundreds keV). The main part of the total GDR width comes from the spreading width $\Gamma^{\downarrow}$ due to coupling to complicated configurations such as $2p2h$ and even more complex ones. The extension of the microscopic descriptions of the width of the g.s. GDR such as those in [2,3] to non-zero temperature ($T \neq 0$) has shown that all these three components of the GDR width depend weakly on temperature [4] - [6]. Meanwhile the experimental systematic has shown that the width of hot GDR increases rapidly at low excitation energies E^* (or temperature T) and saturates at $E^* \geq 130$ MeV in tin isotopes [7] - [13]. The increase of the width at $T \neq 0$ can be described by the thermal fluctuations of nuclear shapes [14-16]. The width of the GDR may depend noticeably on the angular momentum J if the latter reaches a rather high value $J \geq 35\ \hbar$ at $T \simeq 1.5$ - 1.8 MeV in a lighter nucleus ^{106}Sn [17]. Recently we have shown that the coupling of the collective dipole vibration (GDR phonon) to the incoherent particle-particle (pp) and hole-hole (hh) configurations appearing at $T \neq 0$ (the thermal damping), which is de-facto taking shape fluctuations into account, is decisively important for an adequate description of the width's increase and its saturation [18-21]. It has also been concluded that the quantal damping due to coupling to only ph configurations decreases slowly as T increases. The application of this approach in a systematic study of the hot GDR in ^{90}Zr, ^{120}Sn, and ^{208}Pb has shown a reasonable agreement with the experimental data [18,19]. The higher-order graphs such as $1p1h\otimes$phonon, $1p1p\otimes$phonon, $1h1h\otimes$phonon or/and two-phonon ones have not been included explicitly in [18,19], but rather effectively in the parameters of the model. The aim of the present letter is to study the contribution of these higher-oder processes to the width of the hot GDR. For this purpose a complete set of approximate equations will be derived including all the forward-going processes up to two-phonon ones at $T \neq 0$ and applied in numerical calculations in ^{120}Sn and ^{208}Pb. As the present paper is a further development of the approach in [18], the latter will be frequently quoted as I hereafter.

We consider the same model Hamiltonian as in [18,19] for the description of the coupling of collective oscillations (phonons) to the field of ph, pp and hh pairs. This Hamiltonian is composed of three terms:

$$H = \sum_s E_s a_s^\dagger a_s + \sum_q \omega_q Q_q^\dagger Q_q + \sum_{ss'q} F_{ss'}^{(q)} a_s^\dagger a_{s'} (Q_q^\dagger + Q_q). \quad (1)$$

The first term on the right-hand side (RHS) of Eq. 1 is the field of independent single particles, where $a_s^\dagger$ and a_s are creation and destruction operators of a particle or hole state with energy $E_s = \epsilon_s - \epsilon_F$ with ϵ_s being the single-particle energy and ϵ_F - the Fermi surface's energy. We will simply call the energy E_s as the single-particle energy. The second term stands for the phonon field, where $Q_q^\dagger$ and

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Q_q are the creation and destruction operators of a phonon with energy ω_q. The last term describes the coupling between the first two terms. The indices s and s' denote particle (p, $E_p > 0$) or hole (h, $E_h < 0$), while the index q is reserved for the phonon state $q = \{\lambda, i\}$ with multipolarity λ (the projection μ of λ in the phonon index is omitted in the writing for simplicity). The sums in the last two temrs of Eq. (1) are carried over $\lambda \geq 1$.

We introduce the double-time Green functions in the standard notation [22], which describe the following processes: a) *The propagation of a free phonon*: $G_{q',q}(t-t') = \langle\langle Q_{q'}(t); Q_q^\dagger(t')\rangle\rangle$, b) *The transition between a nucleon pair and a phonon*: $\mathcal{G}_{ss',q}(t-t') = \langle\langle a_s^\dagger(t)a_{s'}(t); Q_q^\dagger(t')\rangle\rangle$, c) *The transition between a nucleon pair⊗phonon configuration and a phonon*: $\Gamma^{-;+}_{ss'q',q}(t-t') = \langle\langle a_s^\dagger(t)a_{s'}(t)Q_{q'}(t); Q_q^\dagger(t')\rangle\rangle$, and d) *The transition between two- and one-phonon configurations*: $G^{--;+}_{q_1q_2;q}(t-t') = \langle\langle Q_{q_1}(t)Q_{q_2}(t); Q_q^\dagger(t')\rangle\rangle$. The effect of single-particle damping on the GDR width has been found in I to be rather small up to high T. Therefore we will not consider it here again. The backward-going processes $G^{+;+}_{q';q}(t-t') = \langle\langle Q_{q'}^\dagger(t); Q_q^\dagger(t')\rangle\rangle$, $\Gamma^{+;+}_{ss'q',q}(t-t') = \langle\langle a_s^\dagger(t)a_{s'}^\dagger(t)Q_{q'}^\dagger(t); Q_q^\dagger(t'),\rangle\rangle$, and $G^{++;+}_{q_1q_2}(t-t') = \langle\langle Q_{q_1}^\dagger(t)Q_{q_2}^\dagger(t); Q_q^\dagger(t')\rangle\rangle$ will be omitted as their effects on the damping of the GDR are expected to be negligible. Applying the standard method of the equation of motion for the double-time Green function [22], we obtain a set of coupled equations for an hierarchy of Green functions. Employing the decoupling approximation discussed previously in I, we can close this set to the functions (a) - (d). Making then the Fourier transformation to the energy plane E, we obtain a set of four equations for the Fourier transforms of the Green functions (a) - (d). The first three equations in this set are exact, while the last is approximated due to the decoupling scheme mentioned above. The latter leads to the single-particle occupation number n_s and phonon occupation number ν_q, whose explicit expressions have been derived in I. Eliminating $\mathcal{G}_{ss';q}(E)$ from this set, we obtain one exact equation, which relates $G_{q,q}(E)$ to $\Gamma^{-;+}_{ss',q}(E)$. Eliminating now $G^{--;+}_{q_1q_2;q}(E)$ by expressing it in terms of $\Gamma^{-;+}(E)$, we come to the an approximate equation, which relates functions $\Gamma^{-;+}(E)$ to $\mathcal{G}_{ss',q}(E)$. Making again the decoupling for all the Green functions under the sums in this equation, which truncates the chain at the second order $O[(F_{ss'}^{(q')})^2]$ of the interaction strength $F_{ss'}^{(q')}$, we can express $\Gamma^{-;+}_{ss'q';q}(E)$ in terms of $G_{q',q}(E)$. Inserting $\Gamma^{-;+}_{ss'q',q}(E)$ from the just obtained approximate equation into the exact equation, which relates $G_{q;q}(E)$ to $\Gamma^{-;+}_{ss',q}(E)$, we end up with the final equation for the propagation of a single phonon ($q_1 = q$) in the following form

$$G_{q,q}(E) = \frac{1}{2\pi}[E - \omega_q - P_q(E)]^{-1}. \tag{2}$$

The polarization operator $P_q(E)$ and the vertex function $\mathcal{M}_{ss'}^{q_1q'}(E)$ in Eq. (2) are

$$P_q(E) = \sum_{ss's_1q'} \frac{F_{ss'}^{(q)}}{E - E_{s'} + E_s}\Big[\frac{F_{s's_1}^{(q')}\mathcal{M}_{ss_1}^{qq'}(E)}{E - E_{s_1} + E_s - \omega_{q'}} - \frac{F_{s_1s}^{(q')}\mathcal{M}_{s_1s'}^{qq'}(E)}{E - E_{s'} + E_{s_1} - \omega_{q'}}\Big]. \tag{3}$$

$$\mathcal{M}_{ss'}^{qq'}(E) = \sum_{s_2}\Big\{\frac{(1 - n_{s'} + \nu_{q'})(n_s - n_{s_2})}{E - E_{s_2} + E_s}F_{s's_2}^{(q')}F_{s_2s}^{(q)} - \frac{(n_s + \nu_{q'})(n_{s_2} - n_{s'})}{E - E_{s'} + E_{s_2}}F_{s_2s}^{(q')}F_{s's_2}^{(q)} +$$

$$n_{s_2}(n_s - n_{s'})\Big[\frac{F_{s's}^{(q)}F_{s_2s_2}^{(q')}}{E - \omega_q - \omega_{q'}} + \delta_{qq'}\sum_{q_1}\frac{F_{s's}^{(q_1)}F_{s_2s_2}^{(q_1)}}{E - \omega_{q_1} - \omega_{q'}}\Big]\Big\}. \tag{4}$$

The damping width Γ_{GDR} of the hot GDR located at energy $\omega = \omega_{GDR}$ is defined via the imaginary part of the analytical continuation of the polarization operator $P_q(E)$ into complex energy plane $E = \omega \pm i\varepsilon$ as

$$\Gamma_{GDR} = 2\gamma_q(\omega)|_{\omega=\omega_{GDR}} = 2|\mathrm{Im}P_q(\omega \pm i\varepsilon)|_{\omega=\omega_{GDR}}. \tag{5}$$

The polarization operator $P_q(E)$ includes not only $1p1h$ and $1p1h$⊗phonon processes as in the NFT [3,5] but also $1p1p$, $1h1h$, $1p1p$⊗phonon, $1h1h$⊗phonon and two-phonon ones at the same second order in the interaction strength. In the limit of high T it is easy to see that the vertex function $\mathcal{M}^{qq'}_{ss'}(E)$ in Eqs. (3) and (4) tends to

$$\mathcal{M}^{q_1q'}_{ss'}(E)|_{T\to\infty} \longrightarrow \frac{1}{4}\sum_{s_2}\Big\{\frac{1}{\omega_{q'}}\Big[\frac{E_{s_2}-E_s}{E-E_{s_2}+E_s}F^{(q')}_{s's_2}F^{(q)}_{s_2s} + \frac{E_{s_2}-E_{s'}}{E-E_{s'}+E_{s_2}}F^{(q')}_{s_2s}F^{(q)}_{s's_2}\Big]+$$

$$\frac{1}{2T}(E_{s'}-E_s)\Big[\frac{F^{(q)}_{s's}F^{(q')}_{s_2s_2}}{E-\omega_q-\omega_{q'}} + \delta_{qq'}\sum_{q_1}\frac{F^{(q_1)}_{s's}F^{(q_1)}_{s_2s_2}}{E-\omega_{q_1}-\omega_{q'}}\Big]\Big\}, \qquad (6)$$

which means that it decreases as $O(T^{-1})$ with increasing T because of the factor T^{-1} in front of two-phonon terms on the right-hand side (RHS) of Eq. (6). Neglecting these two-phonon terms would lead to a constant width at high temperature because the first two terms on the RHS of Eq. (6) are independent of T under the assumption that the temperature dependence of the interaction, of the phonon energy, and of the difference $E_{s'} - E_s$ is negligible.

The calculations of the GDR width Γ_{GDR} from Eq. (5) have been carried out in ^{120}Sn and ^{208}Pb within the interval $0 \leq T \leq 6$ MeV. The schematic model in I is employed, according to which the g.s. GDR is generated by a single collective and structureless phonon width energy ω_q closed to the energy E_{GDR} of the g.s. GDR. The single-particle energies, calculated in the Woods-Saxon potentials at $T = 0$ for ^{120}Sn and ^{208}Pb, were used in calculations. The calculations in [16] have shown that the major contribution of the shape fluctuations in the increase of the GDR width at $T \neq 0$ comes from the quadrupole shape fluctuations. Therefore, we retain here only dipole and quadrupole phonons in the two-phonon configuration mixing for simplicity. The parameters of the model have been selected as follows. We first set the ratio $r = F_i^{(2)}/F_i^{(1)}$ ($i = 1, 2$) and choose ω_{q_1} in Eq. (4) to be close to $E_{2_1^+}$. The values of ω_q in Eq. (4) and of $F_1^{(1)}$ are then selected in such a way that the solution $\bar{\omega}$ of the equation for the pole of the Green function in Eq. (2) $\omega - \omega_q - P_q(\omega) = 0$ is equal to the GDR energy $\bar{\omega} = E_{GDR}$ while the total width $\Gamma_{GDR}(\bar{\omega})$ from Eq. (5) reproduces the empirical width of the GDR at $T = 0$. The value of $F_2^{(1)}$ is chosen so that $\bar{\omega}$ is stable while T is varied. The parameters are kept unchanged with changing T. This ensures that all thermal effects come from the microscopic configuration mixings, but not from varying parameters. The calculations have used a value equal to 0.5 MeV for the smearing parameter ε in Eq. (5). The results have been checked to be stable against varying ε within the interval 0.2 MeV$\leq \varepsilon \leq$ 1.0 MeV.

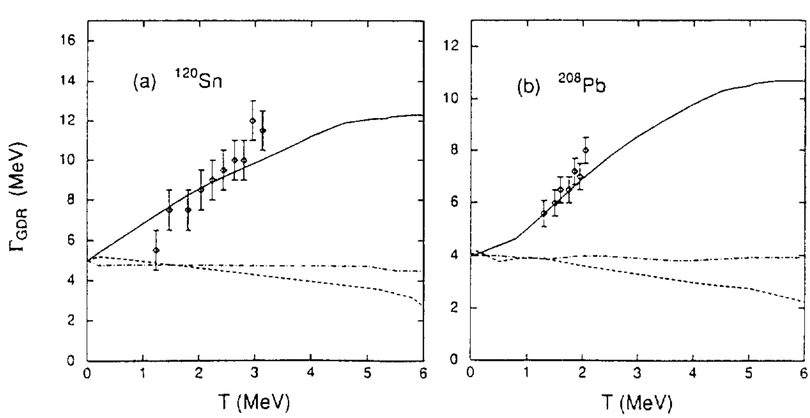

FIG. 1. Width of GDR as a function of temperature for ^{120}Sn (a) and ^{208}Pb (b). The solid curve denotes the total width Γ_{GDR}. The dashed curve denotes the quantal width Γ_Q. The dash-dotted curve stands for the quantal width Γ_Q when the contribution of the two-phonon graphs at $T \neq 0$ is omitted. Experimental data of [13] are shown by open diamonds with error bars.

The total widths Γ_{GDR}, calculated in ^{120}Sn and ^{208}Pb as a function of T, are shown by the solid curves in Fig. 1. An overall agreement is found between theory and the most recent data from the inelastic α-scattering experiments [13]. In both nuclei the region of width's saturation is at around $T \sim 4$ - 6 MeV. In ^{120}Sn the saturated value of the width is about 12 MeV in agreement with the data from [7] - [12]. In ^{208}Pb it is around 10.5 MeV. As has been demonstrated in I, the total width Γ_{GDR} is composed of the quantal width Γ_Q due to coupling of the GDR phonon to ph configurations and the thermal width Γ_T due to coupling to pp and hh configurations at $T \neq 0$. The main conclusion of I is that the behavior of the total width at high temperatures is mostly driven by the thermal width Γ_T since the quantal width Γ_Q decreases as temperature increases. In order to see whether this conclusion still holds within the present more refined approximation we have also switched off the coupling to pp and hh configurations in the sums on the RHS of Eqs. (3) and (4). The results obtained are shown by the dashed curves in Fig. 1. They restore perfectly the results for quantal width Γ_Q in I, which show a clear decrease as T increases. Switching off the two-phonon terms at $T \neq 0$ from Eq. (4) in calculating the width Γ_Q has resulted in a quantal width, which is practically independent of T as shown by the dash-dotted curves in Fig. 1 in agreement with the conclusion within the NFT at $T \neq 0$ [5]. These results show that the NFT indeed includes the graphs, which are most important for an adequate description of the quantal width Γ_Q at $T = 0$, namely the $1p1h$ and $1p1h\otimes$phonon ones. However, in calculating Γ_Q, the NFT neglects entirely the coupling to two-phonon configurations in the vertex function $\mathcal{M}$ in Eq. (4) at $T \neq 0$, which enter in the same second order of the interaction strength $F_{ss'}^{(q)}$. These two-phonon processes at $T \neq 0$ indeed lead to the decrease in the quantal width with increasing T.

In summary we have developed an approach to study the width of the GDR as a function of temperature, which includes all the forward-going processes up to two-phonon ones in the second-order of the interaction strength. The numerical calculations performed in the present paper show that: 1) the total width Γ_{GDR} of the hot GDR arises mainly from the coupling of the GDR vibration to all ph, pp and hh configurations. It increases sharply as temperature increases up to $T \sim 3$ MeV. At higher temperatures the width increase is slowed down to reach a saturated value of around 12 MeV in ^{120}Sn and around 10.5 MeV in ^{208}Pb at $T \sim 4$ - 6 MeV; 2) the quantal width Γ_Q of the GDR due to coupling of GDR vibration to only ph configurations decreases as T increases. Neglecting the two-phonon processes in the expansion to higher-order propagators results in a quantal width, which is almost independent of T.

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