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# Damping of hot giant dipole resonance due to complex configuration mixing

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## Abstract

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  $^{120}\text{Sn}$  and  $^{208}\text{Pb}$ . An overall agreement between theory and experiment is found for the GDR width and shape. 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  $T \sim 4\text{--}6$  MeV. The quantal width  $\Gamma_Q$  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. © 1998 Elsevier Science B.V. All rights reserved.

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The giant dipole resonance (GDR) built on compound nuclear states (hot GDR), has been the subject of a considerable number of experimental and theoretical studies during the last fifteen years (See [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 (phonons) with frequencies around the GDR energy. It can be well described within the random-phase approximation (RPA) and amounts to only a small fraction of the total width. The escape width  $\Gamma^\uparrow$  arises from particle emission. It can be studied 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  $\Gamma^\downarrow$

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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 the three components of the GDR width depend weakly on temperature [4–7]. This is in contradictory to the experimental systematic of the width of hot GDR, which increases rapidly at low excitation energies  $E^*$  (or temperature  $T$ ). At higher excitation energies the observed width increases slowly and even saturates at  $E^* \geq 130$  MeV in tin isotopes [8–14]. The increase of the width at  $T \neq 0$  can be described by the thermal fluctuations of nuclear shapes [15–17]. They were also included in the recent adiabatic coupling model (ACM) [18], whose results agree well with the data from the inelastic  $\alpha$ -scattering experiments in [14] for the GDR width in  $^{120}\text{Sn}$  and  $^{208}\text{Pb}$  within  $1 \text{ MeV} < T \leq 3 \text{ MeV}$ . 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  $T \approx 1.5\text{--}1.8$  MeV in a lighter nucleus  $^{106}\text{Sn}$  [19].

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 [20–22]. 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  $^{90}\text{Zr}$ ,  $^{120}\text{Sn}$ , and  $^{208}\text{Pb}$  has shown a reasonable agreement with the experimental data [20]. 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 [20], but rather effectively in the parameters of the model. The aim of the present letter is to study the contribution of these higher-order 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  $^{120}\text{Sn}$  and  $^{208}\text{Pb}$ . As the present paper is a further development of the approach in [20], the latter will be frequently quoted as I hereafter.

We consider the same model Hamiltonian as in I 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:

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

The first term in the r.h.s. 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  $\epsilon_s$  being the single-particle energy and  $\epsilon_F$  - 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  $Q_q$  are the creation and destruction operators of a phonon with energy  $\omega_q$ . 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  $\lambda$  (the projection  $\mu$  of  $\lambda$  in the phonon index is omitted in the writing for simplicity). The sums in the last two terms in the r.h.s of Eq. (1) are carried over  $\lambda \geq 1$ .

In order to include all the propagators up to two-phonon ones let us introduce the double-time Green functions in the standard notation [23], which describe the following processes: a) *The propagation of a free phonon*:  $G_{q';q}(t-t') = \langle\langle Q_{q'}^\dagger(t); Q_q^\dagger(t') \rangle\rangle$ , 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  $\otimes$  phonon configuration and a phonon*:  $\Gamma_{ss';q}^{\pm,+}(t-t') = \langle\langle a_s^\dagger(t) a_{s'}(t) Q_q^\dagger(t); Q_q^\dagger(t') \rangle\rangle$ , and d) *The transition between two- and one-phonon configurations*:  $G_{q_1 q_2}^{\pm,+}(t-t') = \langle\langle Q_{q_1}^\dagger(t) Q_{q_2}^\dagger(t); Q_q^\dagger(t') \rangle\rangle$ . 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  $G_{q';q}^{+,+}(t-t') = \langle\langle Q_{q'}^\dagger(t); Q_q^\dagger(t') \rangle\rangle$ ,  $\Gamma_{ss';q}^{+,+}(t-t') = \langle\langle a_s^\dagger(t) a_{s'}^\dagger(t) Q_q^\dagger(t); Q_q^\dagger(t') \rangle\rangle$ , and  $G_{q_1 q_2}^{+,+}(t-t') = \langle\langle Q_{q_1}^\dagger(t) Q_{q_2}^\dagger(t); Q_q^\dagger(t') \rangle\rangle$  will be omitted as their effects on the damping of the GDR are expected to be negligible. Applying now the standard method of the equation of motion for the double-time Green function [23], 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) in the form:

$$(E - \omega_q)G_{q;q}(E) - \sum_{s_1 s'_1} F_{s_1 s'_1}^{(q)} \mathcal{G}_{s_1 s'_1; q}(E) = \frac{1}{2\pi}, \quad (2)$$

$$(E - E_{s'} + E_s) \mathcal{G}_{s s'; q}(E) - \sum_{s_1 q'} [F_{s' s_1}^{(q')} \Gamma_{s s_1 q'; q}^{-, +}(E) - F_{s_1 s}^{(q')} \Gamma_{s_1 s' q'; q}^{-, +}(E)] = 0, \quad (3)$$

$$(E - \omega_{q_1} - \omega_{q_2}) G_{q_1 q_2; q}^{-, -; +}(E) - \sum_{s s'} [F_{s s'}^{(q_1)} \Gamma_{s s' q_2; q}^{-, -; +}(E) + F_{s s'}^{(q_2)} \Gamma_{s s' q_1; q}^{-, -; +}(E)] = 0. \quad (4)$$

$$(E - E_{s'} + E_s - \omega_{q'}) \Gamma_{s s' q'; q}^{-, +}(E) - (1 - n_{s'} + \nu_{q'}) \sum_{s_1} F_{s' s_1}^{(q')} \mathcal{G}_{s s_1; q}(E) + (n_s + \nu_{q'}) \sum_{s_1} F_{s_1 s}^{(q')} \mathcal{G}_{s_1 s'; q}(E) - (n_s - n_{s'}) \sum_{q_1} F_{s s'}^{(q_1)} G_{q_1 q'; q}^{-, -; +}(E) = 0, \quad (5)$$

The first three equations in this set are exact, while the last Eq. (5) is approximated due to the decoupling scheme mentioned above. The latter leads to the single-particle occupation number  $n_s$  and phonon occupation number  $\nu_{q'}$ , whose explicit expressions have been derived in I. Eliminating  $\mathcal{G}_{s s'; q}(E)$  from Eqs. (2) and (3), we obtain one exact equation, which relates  $G_{q;q}(E)$  to  $\Gamma_{s s'; q}^{-, +}(E)$ , in the form

$$(E - \omega_q) G_{q;q}(E) - \sum_{s s' s_1 q'} \frac{F_{s s'}^{(q)} [F_{s' s_1}^{(q')} \Gamma_{s' s_1 q'; q}^{-, +}(E) - F_{s_1 s}^{(q')} \Gamma_{s_1 s' q'; q}^{-, +}(E)]}{E - E_{s'} + E_s} = \frac{1}{2\pi}. \quad (6)$$

Eliminating now  $G_{q_1 q_2; q}^{-, -; +}(E)$  by expressing it in terms of  $\Gamma^{-, +}(E)$  using Eq. (4) and inserting the result into Eq. (5), we come to the an approximate equation, which relates functions  $\Gamma^{-, +}(E)$  to  $\mathcal{G}_{s s'; q}(E)$ . Making again the decoupling for all the Green functions under the sums in this equation, which truncates the chain at the second order  $O[(F_{s s'}^{(q')}]^2)$  of the interaction strength  $F_{s s'}^{(q')}$ , we can express  $\Gamma_{s s' q'; q}^{-, +}(E)$  in terms of  $G_{q', q}(E)$  as

$$(E - E_{s'} + E_s - \omega_{q'}) \Gamma_{s s' q'; q}^{-, +}(E) = \sum_{q_1} \mathcal{M}_{s s'}^{q_1 q'}(E) G_{q_1; q}(E). \quad (7)$$

Inserting  $\Gamma_{s s' q'; q}^{-, +}(E)$  from the approximate Eq. (7) into the exact Eq. (6), we end up with the final equation for the propagation of a single phonon ( $q_1 = q$ ) in the following form

$$G_{q, q}(E) = \frac{1}{2\pi} \frac{1}{E - \omega_q - P_q(E)}. \quad (8)$$

The polarization operator  $P_q(E)$  and the vertex function  $\mathcal{M}_{s s'}^{q_1 q'}(E)$  in Eqs. (7) and (8) are

$$P_q(E) = \sum_{s s' s_1 q'} \frac{F_{s s'}^{(q)}}{E - E_{s'} + E_s} \left[ \frac{F_{s' s_1}^{(q')} \mathcal{M}_{s s_1}^{q q'}(E)}{E - E_{s_1} + E_s - \omega_{q'}} - \frac{F_{s_1 s}^{(q')} \mathcal{M}_{s_1 s'}^{q q'}(E)}{E - E_{s'} + E_{s_1} - \omega_{q'}} \right], \quad (9)$$

$$\begin{aligned} \mathcal{M}_{s s'}^{q q'}(E) = \sum_{s_2} \left\{ \frac{(1 - n_{s'} + \nu_{q'}) (n_s - n_{s_2})}{E - E_{s_2} + E_s} F_{s' s_2}^{(q')} F_{s_2 s}^{(q)} - \frac{(n_s + \nu_{q'}) (n_{s_2} - n_{s'})}{E - E_{s'} + E_{s_2}} F_{s_2 s}^{(q')} F_{s' s_2}^{(q)} \right. \\ \left. + n_{s_2} (n_s - n_{s'}) \left[ \frac{F_{s s}^{(q)} F_{s_2 s_2}^{(q')}}{E - \omega_q - \omega_{q'}} + \delta_{q q'} \sum_{q_1} \frac{F_{s s}^{(q_1)} F_{s_2 s_2}^{(q_1)}}{E - \omega_{q_1} - \omega_{q'}} \right] \right\}. \quad (10) \end{aligned}$$

The damping width  $\Gamma_{\text{GDR}}$  of the hot GDR located at energy  $\omega = \omega_{\text{GDR}}$  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

$$\Gamma_{\text{GDR}} = 2\gamma_q(\omega)|_{\omega=\omega_{\text{GDR}}} = 2[\text{Im} P_q(\omega \pm i\varepsilon)]|_{\omega=\omega_{\text{GDR}}}. \quad (11)$$

The polarization operator  $P_q(E)$  includes all  $1s1s'$  and  $1s1s' \otimes$  phonon processes  $((s,s') = (p,h), (p,p')$  and  $(h,h')$ ) and two-phonon ones at the same second order in the interaction strength. In the limit of high  $T$  it is easy to verify that the vertex function  $\mathcal{M}_{s's'}^{qq'}(E)$  in Eqs. (9) and (10) decreases as  $O(T^{-1})$  with increasing  $T$  because of the factor  $T^{-1}$  in front of two-phonon terms on the r.h.s. of Eq. (10). Neglecting these two-phonon terms would lead to a constant width at high temperature because the first two terms on the r.h.s. of Eq. (10) 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 decrease of the quantal width  $\Gamma_Q$  of the GDR as  $O(T^{-1})$  at high  $T$  has also been estimated analytically in the SRPA in [24]. Although we proceeded the calculations up to  $T \leq 6$  MeV, the high  $T$ -limit indicates the decreasing trend of  $\Gamma_Q$  with increasing  $T$  as will be seen below.

The calculations of the GDR width  $\Gamma_{\text{GDR}}$  from Eq. (11) have been carried out in  $^{120}\text{Sn}$  and  $^{208}\text{Pb}$  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  $\omega_q$  closed to the energy  $E_{\text{GDR}}$  of the g.s. GDR. The realistic single-particle energies, calculated in the Woods-Saxon potentials at  $T = 0$  for  $^{120}\text{Sn}$  and  $^{208}\text{Pb}$ , were used in calculations. The calculations in [18] 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. Consequently, from the sums on the r.h.s. of Eqs. (9) and (10) there remain only one dipole phonon with  $q = q'$ , which corresponds to the GDR ( $\lambda = 1$ ) and one quadrupole phonon with  $q_1$ , which corresponds to the energy of  $2_1^+$  state ( $\lambda = 2$ ). 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  $\omega_{q_1}$  in Eq. (10) to be close to  $E_{2_1^+}$ . The values of  $\omega_q$  in Eq. (10) 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. (8)  $\omega - \omega_q - P_q(\omega) = 0$  is equal to the GDR energy  $\bar{\omega} = E_{\text{GDR}}$  while the total width  $\Gamma_{\text{GDR}}(\bar{\omega})$  from Eq. (11) 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. For  $^{120}\text{Sn}$  the set of parameters has been found as:  $\omega_q = 17.0$  MeV,  $F_1^{(1)} = 6.261 \times 10^{-3}$  MeV,  $F_2^{(1)} = 1.845 \times 10^{-1}$  MeV, and  $r = 8.603 \times 10^{-2}$ . For  $^{208}\text{Pb}$  the values  $\omega_q = 13.8$  MeV,  $F_1^{(1)} = 2.7 \times 10^{-3}$  MeV,  $F_2^{(1)} = 9.95 \times 10^{-2}$  MeV, and  $r = 8.8 \times 10^{-1}$  have been selected. These parameters are kept unchanged with changing  $T$  in the present paper. This ensures that all thermal effects come from the microscopic configuration mixings, but not from varying parameters. In fact, as has been shown, e.g. in [5], the first quadrupole state within the RPA has a considerable fragmentation with increasing  $T$ . This offers some possibility to finely tune the parameters of the model, which we will not consider in this stage. The inclusion of higher multiplicities such as octupole phonons in  $^{208}\text{Pb}$ , etc. is also expected to improve the results. However it would certainly make the calculations more complicate. The calculations have used a value equal to 0.5 MeV for the smearing parameter  $\varepsilon$  in Eq. (11). The results have been checked to be stable against varying  $\varepsilon$  within the interval  $0.2 \text{ MeV} \leq \varepsilon \leq 1.0 \text{ MeV}$ .

The total widths  $\Gamma_{\text{GDR}}$ , calculated in  $^{120}\text{Sn}$  and  $^{208}\text{Pb}$  as a function of  $T$ , are shown in Fig. 1 (solid). An overall agreement is found between theory and recent data from the inelastic  $\alpha$ -scattering [14] and heavy-ion fusion experiments [8,10,11]. In both nuclei the region of width's saturation is at around  $T \sim 4-6$  MeV. In  $^{120}\text{Sn}$  the saturated value of the width is about 12 MeV in agreement with the data from [8–13]. In  $^{208}\text{Pb}$  it is around 10.5 MeV. As has been demonstrated in I, the total width  $\Gamma_{\text{GDR}}$  is composed of the quantal width  $\Gamma_Q$  due to coupling of the GDR phonon to  $ph$  configurations and the thermal width  $\Gamma_T$  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  $\Gamma_T$  since the quantal width  $\Gamma_Q$  decreases as temperature increases. In order to see whether this conclusion still holds within the present more refined approximation we have also switched

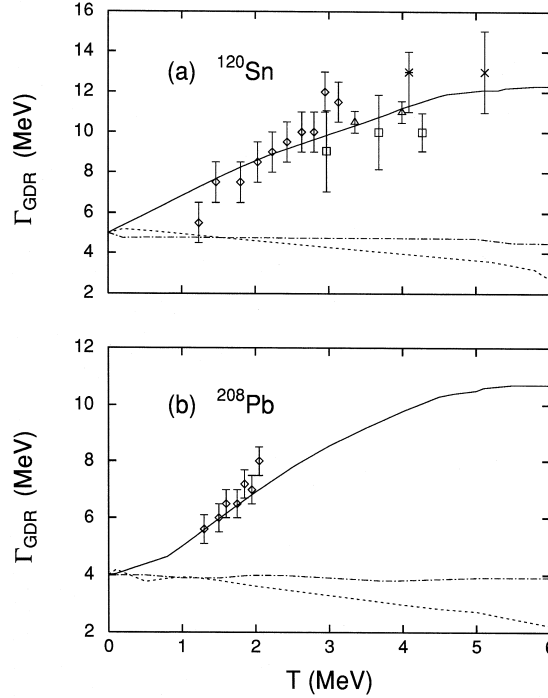


Fig. 1. GDR width as a function of temperature for  $^{120}\text{Sn}$  (a) and  $^{208}\text{Pb}$  (b). *Solid*: total width  $\Gamma_{\text{GDR}}$ ; *dashed*: quantal width  $\Gamma_Q$ ; *dash-dotted*: quantal width  $\Gamma_Q$  when the contribution of the two-phonon graphs at  $T \neq 0$  is omitted. *Diamonds*: data of [14]; *squares* and *triangles*: data by Enders et al. and Hofmann et al. [11], respectively; *cross* and *asterisk*: data from [8] and [10], respectively. A level density parameter  $\sim A/8$  has been used to translate  $E^*$  to  $T$  in fusion data of [8,10]

off the coupling to  $pp$  and  $hh$  configurations in the sums on the r.h.s. of Eqs. (9) and (10). The results obtained are shown by the dashed curves in Fig. 1. They restore perfectly the results for quantal width  $\Gamma_Q$  in I, which show a clear decrease as  $T$  increases. Switching off the two-phonon terms at  $T \neq 0$  from Eq. (10) in calculating the width  $\Gamma_Q$  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  $T \neq 0$  [6,25]. These results show that the NFT indeed includes the graphs, which are most important for an adequate description of the quantal width  $\Gamma_Q$  at  $T = 0$ , namely the  $1p1h$  and  $1p1h \otimes$  phonon ones. However, in calculating  $\Gamma_Q$ , the NFT neglects entirely the coupling to two-phonon configurations 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  $T$  as has been mentioned above. Within the NFT, where the phonon operator has a  $ph$  nature, the exclusion of two-phonon terms was necessary to avoid double counting. We also notice that the NFT solved an RPA equation at  $T \neq 0$  with the sum carried over mainly  $n_h - n_p$  as has been shown by Eq. (24) of [6]. The self-energy graphs within the NFT at  $T = 0$  described just the damping with  $ph$ -vibration doorways as has been shown by Fig. 26 and Eqs. (B8) and (B9) of [3]. The deformation of Fermi surface at  $T \neq 0$  within NFT led to some ground-state correlation graph and single-particle renormalization, whose effect on the GDR damping are known to be small (Figs. 5, 6, and Eqs. (11)–(14) of [6]).

Shown in Fig. 2 are the GDR shapes calculated within our model ((a) and (b)) using the GDR strength function  $S_q^{\text{GDR}}(\omega) = \pi^{-1} \gamma_q(\omega) / [(\omega - \bar{\omega})^2 + (\gamma_q(\omega))^2]$  and those obtained within the ACM [18] ((c) and (d)). The areas of experimental divided spectra as well as the results from the ACM have been taken from [26]. The experimental spectra at low (30–40 MeV) and high (110–120 MeV) excitation energies are put in a one-to-one

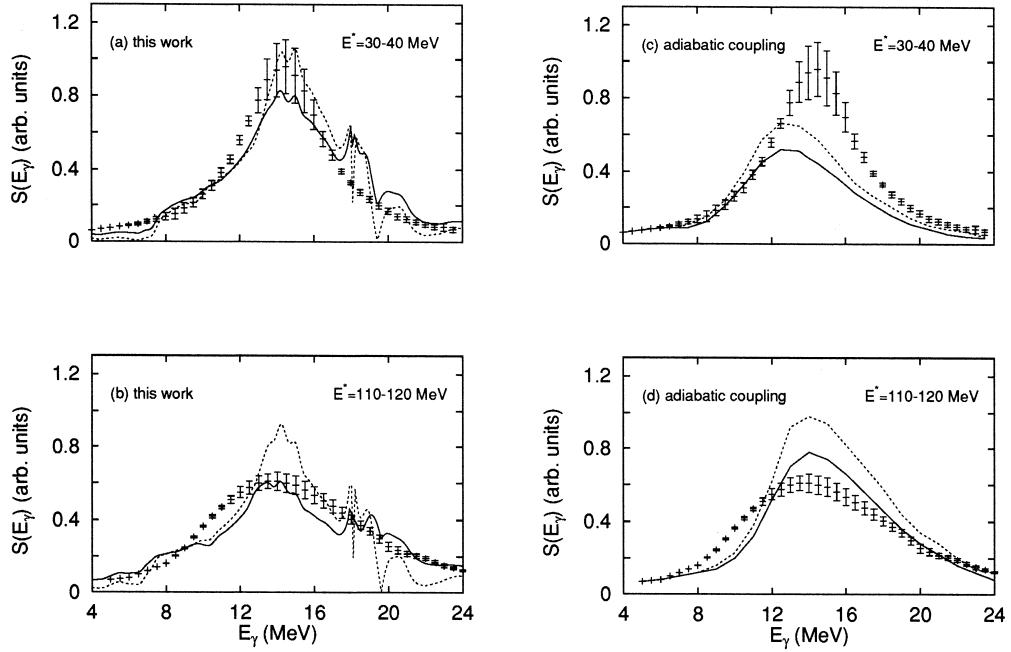


Fig. 2. GDR shapes in  $^{120}\text{Sn}$  at low (30–40 MeV) and high (110–120 MeV) excitation energies. In each panel the vertical bars denote the area of experimental divided spectra. Results of the present work are plotted in (a) and (b), where dashed curves show the shapes calculated without an additional parameter  $\Delta\gamma$  (see the text) while solid curves denote those obtained by adding  $\Delta\gamma$  to  $\gamma_q(\omega)$  ( $\Delta\gamma = 0.5$  MeV in (a) and 2 MeV in (b)). Results of ACM using a strength parameter  $S = 1$  (dashed) and 0.8 (solid) are plotted in (c) and (d).

correspondence with the data of [14] at  $T = 1.24$  MeV and 3.12 MeV, respectively. The calculated shape (dashed curves in Figs. 2(a) and 2(b)) is slightly narrower than the experimental one. Among the possible reasons there are the effects of mixing with more complex configurations, other multipolarities, angular momentum [19], evaporation width [27] at high  $T$ , and coupling to continuum, which are not considered in the present scheme. An effective way to take into account these contributions is to add a smearing parameter  $\Delta\gamma$  to the damping  $\gamma_q(\omega)$  to recover the value of the total width at high  $T$  obtained in I. By doing so the a better agreement between theory and experiment is achieved (solid curves in Figs. 2(a) and 2(b)). It is seen in both cases that our approach offers a reasonable agreement with the observed GDR shape in both regions of low and high excitation energies while the calculated strength function in the ACM does not follow a Lorentzian-like shape at high  $E^*$ . Reducing the value of the GDR energy weighted sum rule value by 20% helped the ACM to improve the agreement with data at high  $E^*$  but worsened it at low excitation  $E^*$ .

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 present paper show that: (1) the total width  $\Gamma_{\text{GDR}}$  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  $^{120}\text{Sn}$  and around 10.5 MeV in  $^{208}\text{Pb}$  at  $T \sim 4$ –6 MeV; (2) the quantal width  $\Gamma_Q$  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  $T$ ; (3) the calculated GDR shape in our model agrees reasonably well with the data. The numerical calculations are limited within a schematic case whose parameters were selected to reproduce the empirical values of the GDR width and its energy at  $T = 0$ . More refined microscopic studies in this direction are highly

desirable since the actual complexity of the wave functions can reveal the detail relationship between the eigenstates of the model Hamiltonian and the representation basis, and, therefore, may shed more light on the physics of the problem. Other ingredients such as the temperature dependence of the single-particle energies, the contribution of the evaporation width, the effects of high values of the angular momentum have been also left out in the present study. While there have been some indications on that these effects may not be crucial within the temperature region under consideration and/or in nuclei with mass number  $A \geq 120$  [11,14,19,26], they certainly deserve thorough studies in the future.

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## References

- [1] K. Snover, *Annu. Rev. Nucl. Part. Sci.* 36 (1986) 545; J.J. Gaardhøje, *Annu. Rev. Nucl. Part. Sci.* 42 (1992) 483.
- [2] V.G. Soloviev, *Theory of Atomic Nuclei – Quasiparticles and Phonons*, IOP, Bristol, 1992.
- [3] G.F. Bertsch, P.F. Bortignon, R.A. Broglia, *Rev. Mod. Phys.* 55 (1983) 287.
- [4] H. Sagawa, G.F. Bertsch, *Phys. Lett. B* 146 (1984) 138.
- [5] N. Dinh Dang, *J. Phys. G* 11 (1985) L125.
- [6] P.F. Bortignon et al., *Nucl. Phys. A* 460 (1986) 149.
- [7] N. Dinh Dang, *Nucl. Phys. A* 504 (1989) 143.
- [8] J.J. Gaardhøje et al., *Phys. Rev. Lett.* 53 (1984) 148; 56 (1986) 1783; 59 (1987) 1409.
- [9] D.R. Chakrabarty et al., *Phys. Rev. C* 36 (1987) 1886.
- [10] A. Bracco et al., *Phys. Rev. Lett.* 62 (1989) 2080; 74 (1995) 3748.
- [11] G. Enders et al., *Phys. Rev. Lett.* 69 (1992) 249; H.J. Hofmann et al., *Nucl. Phys. A* 571 (1994) 301.
- [12] J.H. Le Faou et al., *Phys. Rev. Lett.* 72 (1994) 3321.
- [13] D. Perroutsakou et al., *Nucl. Phys. A* 600 (1996) 131.
- [14] E. Ramakrishnan et al., *Phys. Rev. Lett.* 76 (1996) 2025; T. Baumann et al., *Nucl. Phys. A* 635 (1998) 428.
- [15] M. Gallardo et al., *Nucl. Phys. A* 443 (1985) 415.
- [16] N. Dinh Dang, *J. Phys. G* 16 (1990) 623.
- [17] Y. Alhassid, B. Bush, S. Levit, *Phys. Rev. Lett.* 61 (1988) 1926; Y. Alhassid, B. Bush, *Nucl. Phys. A* 509 (1990) 461; A 531 (1991) 1, 39; B. Bush, Y. Alhassid, *Nucl. Phys. A* 531 (1991) 27.
- [18] W.E. Ormand, P.F. Bortignon, R.A. Broglia, *Phys. Rev. Lett.* 77 (1996) 607; W.E. Ormand et al., *Nucl. Phys. A* 614 (1997) 217.
- [19] M. Matiuzzi et al., *Nucl. Phys. A* 612 (1997) 262.
- [20] N. Dinh Dang, A. Arima, *Phys. Rev. Lett.* 80 (1998) 4145; *Nucl. Phys. A* 636 (1998) 443.
- [21] N. Dinh Dang, F. Sakata, *Phys. Rev. C* 57 (1998) 3032; C 55 (1997) 2872.
- [22] N. Dinh Dang, *Phys. Rep.* 264 (1996) 123.
- [23] D.N. Zubarev, *Sov. Phys. Uspekhi* 3 (1960) 320.
- [24] K. Tanabe, *Nucl. Phys. A* 569 (1994) 27c.
- [25] P. Donati et al., *Phys. Lett. B* 383 (1996) 15.
- [26] G. Gervais, M. Thoennessen, W.E. Ormand, *Phys. Rev. C* 58 (1998) 1377R.
- [27] Ph. Chomaz, *Phys. Lett. B* 347 (1995) 1.