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DAMPING OF ISOVECTOR GIANT DIPOLE RESONANCES IN HOT EVEN-EVEN SPHERICAL NUCLEI

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Abstract: An approach based on the finite temperature quasiparticle phonon nuclear model (FT-QPNM) with the couplings to (2p2h) states at finite temperature taken into account is suggested for calculations of the damping of giant multipole resonances in hot even-even spherical nuclei. The strength functions for the isovector giant dipole resonance (IV-GDR) are calculated in ^{58}Ni and ^{90}Zr for a range of temperatures up to 3 MeV. The results show that the contribution of the interactions with (2p2h) configurations to the IV-GDR spreading width changes weakly with varying temperature. The IV-GDR centroid energy decreases slightly with increasing temperature. The nonvanishing superfluid pairing gap due to thermal fluctuations is included.

1. Introduction

The properties of nuclear collective states in highly excited nuclei have attracted an increasing interest in recent years. Since the first observation of the giant dipole resonances (GDR) in thermally excited nuclei produced from proton (and neutron) capture [1]), in heavy-ion fusion [2-4]) and in heavy-ion deep inelastic reactions [5]) up to now the γ-decay of highly excited nuclei has been measured using large arrays of γ-detectors in many experiments [6-9]).

Many general features of the GDR in hot nuclei with temperature $T \sim 1-3$ MeV and spin $I \leqslant 40\hbar$ are now known [see ref. [10]) for a review]. Generally, with increasing temperature and spin, one observes an enhancement in the broadening of the resonance as well as a downward shift of its centroid energy.

Due to the statistical nature of these processes in highly excited nuclei, many theoretical studies have been performed where the microscopic understanding of the giant resonance damping mechanism is crystallized in the stastical formalism. As a rule, the background of all these microscopic theories is the random phase approximation (RPA) which is generalized to finite temperature (FT-RPA) [11-20]).

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However, these studies have suggested that the RPA, whether temperature dependent or not, involving solely couplings within (1p1h) subspaces (or (1p1p), (1h1h) subspaces additionally at finite temperature) cannot provide a description of the giant resonance damping apart from some discrete splitting [11,13,19]).

As in the case of zero temperature, the giant resonance widths are mainly spreading widths $\Gamma^{\downarrow}$ and result from couplings to more complicated configurations, firstly to (2p2h) nuclear states. The contribution of the escape widths $\Gamma^{\uparrow}$ to the total decay width is connected with the continuum and its effect turns out to be small even at low temperature [17]). In fact, at finite temperature one faces the following two damping mechanisms of the giant resonances in spherical nuclei:

(1) *The damping of multi-quasiparticle states in hot nuclei.* This problem is intimately connected with the damping width of single-particle states in cold nuclei; a question that goes beyond the RPA (or FT-RPA). The answer can be found only by taking into consideration the residual couplings to (2p2h) states or (2p2p), (2h2h) states additionally at finite temperature *T*, see ref. [32]). This is the damping mechanism due to the quantal fluctuations (collision damping).

(2) *The Landau damping that is caused by the statistical fluctuations.* In the FT-RPA treated at the level of mean field theory, the Landau damping turns out to be unimportant damping mechanism or one that does not vary much with temperature [15–17,19]). What role the Landau damping plays when quantal fluctuations due to the couplings to (2p2h) configurations from (1), are taken into account at finite temperature, is however, still an open question.

As the theoretical understanding of damping of nuclear collective motion is still not clear, the damping mechanism has received much attention in the last years [21–17]) and the two above-mentioned questions still await quantitative answers.

In hot and strongly rotating nuclei the shape deformations or high spin that split the GDR into several components give an important contribution to the broadening of the giant resonance width. This question has been investigated in detail in many papers within the FT-RPA framework [12,23,25,29,30]) and we do not touch upon it here.

Recently, we have extended the quasiparticle-phonon nuclear model (QPNM) [see ref. [31]) and references therein], which has a successful application in the description of many features of collective excitations, including the damping width of giant resonances [22]) in cold nuclei to the case of finite temperature [19,32,33]). Using the Green function technique at finite temperature we have obtained in refs. [32,33]) an explicit set of basic equations for the finite temperature QPNM (FT-QPNM) including also the effects of phonon correlations in the ground state and phonon scattering. We have also presented in refs. [32,33]) a diagrammatic illustration for the physical processes in the derived equations, and pointed out the new diagrams that appear exceptionally at finite temperature [cf. also refs. [34,35])]. Some time later an alternative approach, taking into account the residual couplings to 2p2h states at finite temperature, was suggested by Yannouleas and Jang [36]) as a generalization of their second RPA formalism [37]) to finite temperature. Using the finite temperature

Matsubara formalism Bortignon *et al.* [24]) have recently calculated the damping of single-particle and vibrational motions to lowest order in the coupling between the particles and vibrations ("1p1h⊕phonon") based on the approach of the nuclear field theory (NFT) [21]). Their results have shown that the positions and widths of the giant resonances in even-even spherical nuclei are rather stable with varying temperature. Although lately there have been some theoretical efforts to consider the coupling of the giant vibration to the thermal fluctuations of the nuclear surface [23,25,26,38]), the full microscopic quantitative calculations including the interactions with (2p2h) configurations and statistical fluctuations at finite temperature are still absent up to now.

In the present work we therefore follow the microscopic approach of the above-mentioned FT-QPNM suggested by us in refs. [19,32,33]) to the damping of the giant resonances in even-even spherical nuclei at finite temperature. In our approach the model hamiltonian is expressed in terms of the thermal phonon (two-quasiparticle) operators whose structure is calculated in the FT-RPA [19]). The couplings to the (2p2h) configurations are considered by introducing the wave functions of excited states consisting of one- and two-phonon components at finite temperature by analogy with the zero temperature QPNM [31]). We also emphasize the interplay between the Landau damping due to statistical (or thermodynamical) fluctuations and quantal fluctuations (collision damping or the interactions with (2p2h) configurations), which is quite important for describing properly the observed width of giant resonances at finite temperature [19,28]). Although the calculations in the present work are performed solely for the IV-GDR at finite temperature, for which a great number of experimental data have been collected in hot nuclei, the investigation of higher multipole resonances in the framework of the FT-QPNM is straightforward. They will be displayed in our forthcoming publications.

It is noteworthy that our approach includes the superfluid pairing correlations at finite temperature in a consequent way. It has been shown in a series of works [28,39-42]) that at finite temperature, in virtue of large thermal fluctuations due to the finiteness of nuclei, the pairing gap does not vanish and remains sufficiently large even at moderate temperatures (up to 3-4 MeV). Therefore, at the temperature of interest its effect cannot be neglected. The influence of the pairing correlations at finite temperature on the strength distribution for collective excitations has been investigated in the FT-RPA in many studies [19,20,28,41,42]). It has also been included in the investigations of fluctuations in selected observables in several works [39,43,44]). In our approach the superfluid pairing correlations are included in the definition of the Bogolubov quasiparticles at finite temperature. Finally, the microscopic calculations taking the couplings to (2p2h) states into consideration at finite temperature in the framework of our approach can provide a reappraisal for the semiclassical description of IV-GDR in hot nuclei. The latter has recently been proposed in the extended Vlasov approach by Cai and Di Toro and apparently gives a width increasing moderately with temperature [28]).

The present paper is organized as follows:

Sect. 2 summarizes the formalism of the FT-QPNM, a generalization of the QPNM [31]) to finite temperature, that has been suggested by us in recent publications [19,32]) and in ref. [33]). The formulae for the damping derived in this approach are presented in brief. Even though these formulae are quite complicated in structure, they will be cast here in a form showing a simply physical interpretation for the most important terms.

Sect. 3 establishes briefly the connection between the FT-QPNM equations and the method of strength functions that we have extended to finite temperature for calculating the centroid energies and spreading widths of giant resonances in hot nuclei. We also discuss in this section the choice of parameters we use in our approach for numerical calculations at finite temperature.

Sect. 4 is devoted to the numerical calculated of the strength functions for the GDR in two even-even spherical nuclei ^{58}Ni and ^{90}Zr for a range of temperatures. The contributions of the Landau damping and the quantal fluctuations are considered and discussed. The relation with the recent experimental data is pointed out as well.

Sect. 5 provides a summary and some conclusions.

2. Formalism of the model

The FT-QPNM uses the model hamiltonian of the QPNM [31]). For the investigations of electric transitions (Eλ-transitions) with multipolarity λ (for GDR $\lambda = 1$) we adopt the hamiltonian which consists of the terms describing, respectively the nucleon motion in the mean-field H_{av}, the monopole superfluid pairing interactions H_{pair} and the residual interactions in the form of separable multipole (ph) forces H_{ph}. The effect of the spin-multipole interactions on the generation of Eλ states has been studied in detail within the QPNM in ref [45]) for cold nuclei. It turns out to be negligible in the strength distribution of one-phonon states. We therefore neglect the contribution of spin-multipole forces in our consideration. The explicit form of this hamiltonian has already been given and discussed in detail in the QPNM elsewhere [see e.g. ref. [31])] and we do not repeat it here. In the following subsections we represent the scheme of our approach incorporating one- and two-phonon states as well as the superfluid pairing interactions at finite temperature.

2.1. THE SUPERFLUID PAIRING GAP AT FINITE TEMPERATURE

By using the Bogolubov canonical transformation one can express the abovementioned FT-QPNM hamiltonian in terms of the quasiparticle operators [31]). In comparison with the zero temperature case, the difference at finite temperature consists, first, in the definition of these quasiparticles. In fact, at finite temperature one has to solve the temperature dependent BCS equations (FT-BCS) to find the quasiparticle energy ε_j which equals [11,46])

$$\varepsilon_j(T) = \sqrt{(E_j - \lambda)^2 + \Delta_T^2} \tag{1}$$

where Δ_T is defined by the FT-BCS equation

$$\Delta_T = G \sum_j (j+\tfrac{1}{2}) u_j v_j (1-2n_j) \tag{2}$$

with n_j being the quasiparticle number occupation at temperature T:

$$n_j = [\exp(\varepsilon_j/T)+1]^{-1}, \tag{3}$$

u_j, v_j are the Bogolubov coefficients and G is the pairing constant. The single-particle energies E_j and the chemical potential λ are found at finite temperature T by the average number conserving condition

$$N = \sum_j (j+\tfrac{1}{2}) \left[1 - \frac{E_j - \lambda}{\varepsilon_j}(1-2n_j)\right] \tag{4}$$

and depend, in general, also on T. It is well-known in the FT-BCS theory that the gap Δ_T from (2) collapses at a critical temperature $T_{\rm crit}$ which equals $T_{\rm crit} = 0.567 \times \Delta_{T=0}$ [refs. [11,46])] with the smooth spacing energy spectrum or is of an order of $\frac{1}{2}\Delta_{T=0}$ because the FT-BCS formalism is presented without exact number projection. However, by taking the thermal fluctuations into account the pairing gap does not vanish at finite temperature and the phase transition from superfluid to normal is completely smeared out [38–42]). The thermal average pairing gap is defined as [38–42,46])

$$\langle \Delta_T \rangle = \int_0^\infty \Delta_T P(\Delta_T, T)\, \mathrm{d}\Delta_T \Big/ \int_0^\infty P(\Delta_T, T)\, \mathrm{d}\Delta_T, \tag{5}$$

where $P(\Delta_T, T)$ is the probability that the pairing gap takes any given value Δ_T at temperature T

$$P(\Delta_T, T) = B(\Delta_T) \exp[-F(\Delta_T)/T]. \tag{6}$$

In eq. (6) $F(\Delta_T)$ is the free energy of the system consisting of the nuclear mean-field with the superfluid pairing correlations [11,19,41]). $B(\Delta_T)$ is the Δ-dependence of the mass [40]). In general, this method can be improved by taking into account the particle number fluctuations due to the number nonconserving problem in the Bogolubov transformation. There have been several approximate approaches to resolve this problem at zero temperature [47,48]). One of them has recently been extended by us to finite temperature [41,42]). However, as has also been discussed in ref. [47]) for the zero temperature case and in ref. [44]) for finite temperature the particle number fluctuations can always be considered effectively by some slight renormalization of the pairing constant G in eq. (2) which must be chosen in our calculations from the experimental pairing energies at $T=0$ [cf. e.g. ref. [31])]. Therefore, for simplicity, we shall not consider the contribution of the particle number fluctuations in such a renormalization in the present work. The detailed studies of the effect of nonvanishing gap (5) on many nuclear characteristics at finite temperature and high spin have been performed in many papers to which we refer the readers for further consideration [see e.g. refs. [39-42,44]]

2.2. THE THERMAL ONE-PHONON STATES IN THE FT-RPA

After defining the quasiparticles at finite temperature we introduce the thermal phonon operators [19,49]

$$
\begin{aligned}
Q^{+}_{\lambda\mu i}(T) &= \frac{1}{2}\sum_{jj'}\{\tilde{\psi}^{\lambda i}_{jj'}A^{+}_{\lambda\mu}(jj') - \tilde{\varphi}^{\lambda i}_{jj'}A_{\lambda\tilde{\mu}}(jj') + \tilde{\xi}^{\lambda i}_{jj'}B_{\lambda\mu}(jj') - \tilde{\zeta}^{\lambda i}_{jj'}B^{+}_{\lambda\tilde{\mu}}(jj')\}\,,\\
Q_{\lambda\mu i}(T) &= [Q^{+}_{\lambda\mu i}(T)]^{+}\,.
\end{aligned} \tag{7}
$$

The operators A^{+}, A, B, B^{+} are the well-known two-quasi-particle operators

$$
\begin{aligned}
A^{+}_{\lambda\mu}(jj') &= \sum_{jj'}\langle jmj'm'|\lambda\mu\rangle\alpha^{+}_{jm}\alpha^{+}_{j'm'} \equiv [\alpha^{+}_{j}\otimes\alpha^{+}_{j'}]_{\lambda\mu}\,,\\
B_{\lambda\mu}(jj') &= -\sum_{jj'}\langle jmj'm'|\lambda\mu\rangle(-)^{j'-m'}\alpha^{+}_{jm}\alpha_{j'-m'} \equiv -[\alpha^{+}_{j}\otimes\tilde{\alpha}_{j'}]_{\lambda\mu}\,,\\
A = (A^{+})^{+}\,, \qquad B &= (B^{+})^{+}\,, \qquad A_{\lambda\tilde{\mu}}(jj') \equiv (-)^{\lambda-\mu}A_{\lambda-\mu}(jj')\,.
\end{aligned} \tag{8}
$$

In view of the last two terms containing B and B^{+} in eq. (7), the thermal phonon operators $Q^{+}(T)$ and $Q(T)$ are different from the conventional microscopic RPA phonon operators employed at zero temperature in the QPNM [31] when $\langle B\rangle \sim n_j$ is zero.

Expressing the FT-QPNM hamiltonian in terms of the thermal phonon operators (7) and applying the linearization method for equations of motion, we have obtained in ref. [19] the set of the FT-RPA equations for finding the energies $\omega_{\lambda i}(T)$ of the thermal one-phonon excitations at temperature T and the phonon amplitudes $\tilde{\psi}$, $\tilde{\varphi}$, $\tilde{\xi}$, $\tilde{\zeta}$ from eq. (7). Such a system of the FT-RPA equations has also been derived by other authors based on somewhat different methods [11,12]. Since these materials are now well-known and have already been published we refer the readers to refs. [11,12,19] for a detailed consideration. The expressions for the phonon amplitudes $\tilde{\psi}$, $\tilde{\varphi}$, $\tilde{\xi}$, $\tilde{\zeta}$ are found to be

$$
\begin{aligned}
\tilde{\psi}^{\lambda i}_{jj'} &= \frac{1}{\sqrt{2\mathcal{Y}_{\lambda i}(T)}}\,\frac{f^{(\lambda)}_{jj'}u^{(+)}_{jj'}}{\varepsilon_j+\varepsilon_{j'}-\omega_{\lambda i}(T)}\sqrt{1-n_j-n_{j'}}\,,\\
\tilde{\varphi}^{\lambda i}_{jj'} &= \frac{1}{\sqrt{2\mathcal{Y}_{\lambda i}(T)}}\,\frac{f^{(\lambda)}_{jj'}u^{(+)}_{jj'}}{\varepsilon_j+\varepsilon_{j'}+\omega_{\lambda i}(T)}\sqrt{1-n_j-n_{j'}}\,,\\
\tilde{\xi}^{\lambda i}_{jj'} &= -\frac{1}{\sqrt{2\mathcal{Y}_{\lambda i}(T)}}\,\frac{f^{(\lambda)}_{jj'}v^{(-)}_{jj'}}{\varepsilon_j-\varepsilon_{j'}-\omega_{\lambda i}(T)}\sqrt{n_{j'}-n_j}\,, \qquad (\varepsilon_j>\varepsilon_{j'})\\
\tilde{\zeta}^{\lambda i}_{jj'} &= -\frac{1}{\sqrt{2\mathcal{Y}_{\lambda i}(T)}}\,\frac{f^{(\lambda)}_{jj'}v^{(-)}_{jj'}}{\varepsilon_j-\varepsilon_{j'}+\omega_{\lambda i}(T)}\sqrt{n_{j'}-n_j}\,, \qquad (\varepsilon_j>\varepsilon_{j'})
\end{aligned} \tag{9}
$$

where the norm $\mathcal{Y}_{\lambda i}(T)$ has been given in ref. [19]) with $f_{jj'}^{(\lambda)} \equiv \langle j' \| iR_\lambda(r) Y_\lambda \| j \rangle$ being the single-particle matrix elements corresponding to the separable multipole interactions; $u_{jj'}^{(+)} = u_j v_{j'} + u_{j'} v_j$, $v_{jj'}^{(-)} = u_j u_{j'} - v_j v_{j'}$. Thus, in this way the (1p1h) configurations at finite temperature have been interpreted in terms of thermal one-phonon excitations in the FT-RPA. We note that, as has been discussed in detail in refs. [11,12,19]) at finite temperature due to the presence of the operators B and B^+ in eq. (7) the (pp) and (hh) transitions appears besides the usual (ph) transitions. As has been considered in detail in refs. [11-13,19]) these new configurations lead to a new discrete splitting of one-phonon states. In fact, under their influence some low-lying strongly collective states can spread over several new states with less collectivity [13,19]). As concerns high lying states in the resonance region, the new states correspond, as a rule, to very small reduced transition probabilities $B(\mathrm{E}\lambda\uparrow, \omega_{\lambda i}(T))$ in the temperature region of interest ($T \leqslant 5$ MeV) [19]). Moreover, the Landau splitting calculated in the FT-RPA with including these (pp), (hh) configurations changes not much at finite temperature [13,19]) as has been mentioned in the Introduction. Nevertheless, these new configurations complicate the thermal one-phonons space. Therefore, in the study of the couplings to (2p2h) configurations, they should undoubtedly be taken into account, as has been done in the present work making use of eq. (7).

The inverse transformations for operators A^+, A, B and B^+ can be obtained from the thermal phonon operators $Q^+_{\lambda\mu i}(T)$ and $Q_{\lambda\mu i}(T)$ (7) and their amplitudes (9) in the FT-RPA as

$$\begin{aligned}
A^+_{\lambda\nu}(j_1 j_2) &= \sum_i \{\tilde{\psi}^{\lambda i}_{j_1 j_2} Q^+_{\lambda\mu i}(T) + \tilde{\varphi}^{\lambda i}_{j_1 j_2} Q_{\lambda\tilde{\mu}_i}(T)\}(1 - n_{j_1} - n_{j_2})^{1/2} \\
A_{\lambda\mu}(j_1 j_2) &= \sum_i \{\tilde{\psi}^{\lambda i}_{j_1 j_2} Q_{\lambda\mu i}(T) + \tilde{\varphi}^{\lambda i}_{j_1 j_2} Q^+_{\lambda\tilde{\mu} i}(T)\}(1 - n_{j_1} - n_{j_2})^{1/2} \\
B_{\lambda\mu}(j_1 j_2) &= \sum_i \{\tilde{\xi}^{\lambda i}_{j_1 j_2} Q^+_{\lambda\mu i}(T) - \tilde{\zeta}^{\lambda i}_{j_1 j_2} Q_{\lambda\tilde{\mu} i}(T)\}(n_{j_2} - n_{j_1})^{1/2} \\
B^+_{\lambda\mu}(j_1 j_2) &= \sum_i \{\tilde{\xi}^{\lambda i}_{j_1 j_2} Q_{\lambda\mu i}(T) - \tilde{\zeta}^{\lambda i}_{j_1 j_2} Q^+_{\lambda\tilde{\mu} i}(T)\}(n_{j_2} - n_{j_1})^{1/2} .
\end{aligned} \tag{10}$$

Eqs. (10) are useful in expressing the operators in the FT-QPNM hamiltonian through the thermal one-phonon operators (7).

The FT-RPA excitations (the thermal one-phonon excitations) are now described by the wave functions

$$|\Psi^{(1)}_{JMI}\rangle_T = Q^+_{JMI}(T)|0, T\rangle , \tag{11}$$

where $|0, T\rangle$ is the thermal ground-state wave function which is taken to be the thermal phonon vacuum in even-even spherical nuclei:

$$Q_{JMI}(T)|0, T\rangle = \langle 0, T|Q^+_{JMI}(T) = 0 . \tag{12}$$

If one wishes to apply the variational procedure at finite temperature to obtain the above-mentioned FT-RPA equations based on the wave functions (11), one must have in mind that the average over the thermal vacuum (11) is identical to the average over the canonical ensemble [19,32,33]) with the hamiltonian H

$$\langle\cdots\rangle \equiv \mathrm{Sp}\,[\ldots \exp(-H/T)]/\mathrm{Sp}\,[\exp(-H/T)]\,. \tag{13}$$

2.3. THE COUPLINGS TO THE THERMAL TWO-PHONON STATES

In order to take into account the couplings to (2p2h) configurations we follow the analogous way in the zero temperature QPNM. Namely, we construct the wave functions of excited states, which consist of the thermal one- and two-phonon components [32,33])

$$|\Psi^{(1-2)}_{JMI}\rangle_T = \Omega^+_{JMI}(T)|0, T\rangle\,, \tag{14}$$

where

$$\Omega^+_{JMI}(T) = \Big\{\sum_{\nu} R^I_\nu(J) Q^+_{JM\nu}(T) + \sum_{\substack{\lambda_1 i_1 \\ \lambda_2 i_2}} P^{\lambda_2 i_2}_{\lambda_1 i_1}(JI)[Q^+_{\lambda_1 i_1}(T) \otimes Q^+_{\lambda_2 i_2}(T)]_{JM}$$
$$+ \sum_{\substack{\lambda_1 i_1 \\ \lambda_2 i_2}} S^{\lambda_2 i_2}_{\lambda_1 i_1}(JI)[Q^+_{\lambda_1 i_1}(T) \otimes \tilde{Q}_{\lambda_2 i_2}(T)]_{JM}\Big\} \tag{15}$$

with the notation used in eq. (8). The last terms in eq. (15) appear only at finite temperature and lead to the so-called phonon scattering effect [32,33]).

In ref. [32]) using the wave functions similar to eqs. (14), (15) whose coefficients R, P, S can be put in the one-to-one correspondence with the defined temperature Green functions, and by taking into account the backward one- and two-phonon amplitudes [51]) as well, we have derived explicitly a system of approximate basic equations in the framework of the FT-QPNM including the couplings to the forward two-phonon configurations, the phonon ground-state correlations [51]) and the phonon scattering at finite temperature [32]) [cf. also ref. [35])]. The diagrams describing these physical processes have also been displayed and discussed in detail [32,33]). This set of equations is very complicated and in the recent level of the computation technique it is hardly convenient for numerical calculations in realistic nuclei. Therefore, for an application of these equations to the calculations of the giant resonance damping in hot nuclei this set must be simplified. The effect of phonon ground state correlations has been evaluated in the QPNM using an oversimplified two level model in ref. [50]). In ref. [32]) a similar schematic model for the heated phonon gas with the phonon structure and energies calculated from the RPA has been considered to evaluate the effects of phonon ground state correlations and phonon scattering at finite temperature. These schematic evaluations have shown that at the temperatures of interest the effects of phonon ground-state correlations and phonon scattering can be expected to be small as compared to the effect of two-phonon creation or annihilation (the terms $\sim[Q^+\otimes Q^+]$ in eq. (15)). On the other hand, the Pauli principle can be taken into account already in an approximate way by analogy with the zero temperature case [51]). Our calculations in the zero temperature QPNM have shown that in the damping of the giant resonances two procedures taking the Pauli principle into consideration approximately and exactly [34]) give nearly the same results in even-even spherical nuclei [22]). Furthermore, the approximate procedure requires much shorter computer time. Thus, neglecting the phonon ground-state correlations and the phonon scattering effects at finite temperature, we can pronouncedly simplify the set of basic FT-QPNM equations obtained in refs. [32,33]) to retain only the couplings to the forward two-phonon configurations. This can be done by putting all the coefficients ζ, V and W in eqs. (21) or ref. [32]) to be zero*. In the sense of the operators of excited states (14) it is equivalent to neglecting the last terms in eq. (15) containing $[Q^+(T)\otimes\tilde{Q}(T)]$ and corresponding to the phonon scattering effect [in eq. (15) the effect of phonon ground-state correlations due to the backward terms $\sim Q(T)$ and $[Q(T)\otimes Q(T)]$ has not been taken into account]. After some transformations, by analogy with the zero temperature case and having in mind the average (13), we obtain from the simplified systems of FT-QPNM equations the secular equations for finding energies η of excited states at finite temperature in the form

$$\det\left\|[\omega_{Ji}(T)-\eta]\delta_{ii'}-\frac{1}{2}\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}}\frac{U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)_T U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji')_T(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2})}{\omega_{\lambda_1 i_1}(T)+\omega_{\lambda_2 i_2}(T)-\eta}\right\|=0\,, \tag{16}$$

where $\nu_{\lambda i}$ is the phonon occupation number at temperature T

$$\nu_{\lambda i}=[\exp(\omega_{\lambda i}(T)/T)-1]^{-1} \tag{17}$$

on the thermal one-phonon level with energy $\omega_{\lambda i}(T)$ calculated in the FT-RPA [19]). The terms containing coefficients $U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)_T$ appear as manifestation of the anharmonic effects due to the quasiparticle-phonon interactions [31-33]). They govern the couplings to (2p2h) configurations in the above-mentioned approximation without phonon ground-state correlations and phonon scattering at finite temperature. In our approach the first-order terms containing B^+B (eq. (8)) is included in the FT-RPA [19]) while the higher orders associated with the couplings configurations more complicated than (2p2h) are neglected. The explicit form of the thermal coefficients $U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)_T$ can be obtained by using the inverse transformations (10) to express the expansion of the operator $B(B^+)$ in the order of $\sim A^+A$ through the thermal phonon operators $Q^+(T)$ and $Q(T)$ (eq. (7)). After that one obtains the approximated FT-QPNM hamiltonian in terms of the thermal phonon operators $Q^+(T)$ and $Q(T)$ in the same forms as in the zero temperature case [32-34,50,52]). Another way to derive the coefficients $U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)_T$ at finite temperature is to replace the amplitudes ψ, φ in the expression for $U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)$ at zero temperature [32-34,50,52])

* We apologize for using in eq. (9) and ref. [32]) the same letter ζ to label different quantities.

respectively, by

$$(\psi_{jj'}^{\lambda i})_{T=0} \rightarrow \Theta_{jj'}^{\lambda i} \equiv \begin{pmatrix} \tilde{\psi}_{jj'}^{\lambda i} \\ \tilde{\xi}_{jj'}^{\lambda i} \end{pmatrix}, \qquad (\varphi_{jj'}^{\lambda i})_{T=0} \rightarrow \Xi_{jj'}^{\lambda i} \equiv \begin{pmatrix} \tilde{\varphi}_{jj'}^{\lambda i} \\ \tilde{\zeta}_{jj'}^{\lambda i} \end{pmatrix}. \tag{18}$$

Employing now the symmetry properties of the amplitudes $\tilde{\xi}$ and $\tilde{\zeta}$ from eq. (9), we find that the terms containing the products of the type $\sim\tilde{\xi}\xi$, $\tilde{\xi}\tilde{\xi}$ and $\tilde{\zeta}\tilde{\zeta}$ are completely cancelled in the expressions for $U_{\lambda_1 i_1}^{\lambda_2 i_2}(Ji)$ at $T \neq 0$ to retain only the terms containing $\tilde{\psi}\tilde{\psi}$, $\tilde{\psi}\tilde{\varphi}$ and $\tilde{\varphi}\tilde{\varphi}$. Thus, we obtain for the coefficients $U_{\lambda_1 i_1}^{\lambda_2 i_2}(Ji)_T$ at finite temperature the same expressions as for tthose at zero temperature [32-34,50,52])

$$\begin{aligned} U_{\lambda_1 i_1}^{\lambda_2 i_2}(Ji)_T &= (-)^{\lambda_1+\lambda_2-J}\sqrt{(2\lambda_1+1)(2\lambda_2+1)} \\ &\times \sum_{j_1 j_2 j_3} \left\{ \frac{f_{j_1 j_2}^{(J)}}{\sqrt{2\mathcal{Y}_{Ji}(T)}} v_{j_1 j_2}^{(-)} \begin{Bmatrix} \lambda_1 & \lambda_2 & J \\ j_2 & j_1 & j_3 \end{Bmatrix} (\tilde{\psi}_{j_3 j_1}^{\lambda_1 i_1} \tilde{\varphi}_{j_2 j_3}^{\lambda_2 i_2} + \tilde{\psi}_{j_2 j_3}^{\lambda_2 i_2} \tilde{\varphi}_{j_3 j_1}^{\lambda_1 i_1}) \right. \\ &+ \frac{f_{j_1 j_2}^{(\lambda_1)}}{\sqrt{2\mathcal{Y}_{\lambda_1 i_1}(T)}} v_{j_1 j_2}^{(-)} \begin{Bmatrix} \lambda_1 & \lambda_2 & J \\ j_3 & j_2 & j_1 \end{Bmatrix} (\tilde{\psi}_{j_3 j_1}^{\lambda_2 i_2} \tilde{\psi}_{j_2 j_3}^{Ji} + \tilde{\varphi}_{j_2 j_3}^{Ji} \tilde{\varphi}_{j_3 j_1}^{\lambda_2 i_2}) \\ &+ \left. \frac{f_{j_1 j_2}^{(\lambda_2)}}{\sqrt{2\mathcal{Y}_{\lambda_2 i_2}(T)}} v_{j_1 j_2}^{(-)} \begin{Bmatrix} \lambda_1 & \lambda_2 & J \\ j_1 & j_3 & j_2 \end{Bmatrix} (\tilde{\psi}_{j_3 j_1}^{\lambda_1 i_1} \tilde{\psi}_{j_2 j_3}^{\lambda_2 i_2} + \tilde{\varphi}_{j_2 j_3}^{\lambda_2 i_2} \tilde{\varphi}_{j_3 j_1}^{\lambda_1 i_1}) \right\}. \end{aligned} \tag{19}$$

However, although there is a formal similarity in the expression for $U_{\lambda_1 i_1}^{\lambda_2 i_2}(Ji)$ at zero and finite temperatures, the difference in the two cases is radical. Indeed, the amplitudes $\tilde{\psi}$, $\tilde{\varphi}$ and the norm $\mathcal{Y}(T)$ in eq. (19) are calculated from the FR-RPA equations [19]) (eqs. (9)) whose structure is drastically different from the one of the RPA equations. The coefficients $v_{j_1 j_2}^{(-)}$ in eq. (19) are defined by solving the FT-BCS equations and therefore depend also upon temperature T. Eqs. (16) are the central equations taking into account the interactions with (2p2h) configurations at finite temperature in our approach.

In the diagrammatic representation of the QPNM (FT-QPNM) [31-35]) eqs. (16) correspond to the graph in fig. 1a. The phonon ground-state correlations and phonon scattering, neglected in the present work, are illustrated respectively by the graphs 1b and 1c. There are also mutual combinations between the vertices of graphs 1a–c which we do not depict here. It has been shown in refs. [22,31,34]) that by replacing one intermediate noncollective phonon in these graphs by two quasiparticles, one obtains from them the graphs of the NFT. In fig. 1 this transition is denoted by the arrows and the corresponding NFT graphs are depicted in fig. 1d–f. Therefore, eqs. (16) incorporate the NFT graphs given in fig. 1d which are the most important graphs [24]). In ref. [24]), the results obtained in the Matsubara Green function formalism correspond to the sum of all time orderings of the NFT graphs in fig. 1d–f. However, as has been discussed by Bortingnon *et al.* in ref. [24]), the graphs 1e, f will be important

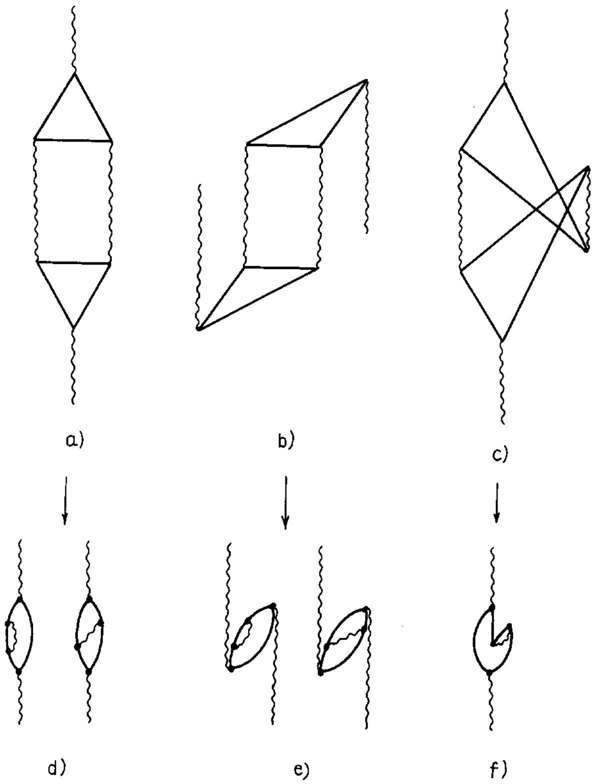

Fig. 1. Basic graphs of the FT-QPNM; (a) the two-phonon creation or annihilation process; (b) the phonon ground-state correlations; (c) the phonon scattering. The graphs (d)–(f) are the graphs of the NFT at finite temperature, which can be obtained from the FT-QPNM graphs (a)–(c) as discussed in text.

only if the vibrational states are strongly collective in the sense that the backward-going amplitudes of the RPA solution are significant. As in the giant resonance region these amplitudes are small as compared to the forward-going ones, the contribution of the NFT graphs in fig. 1e, f can be expected to be negligible at moderate temperatures. For a more detailed discussion of the NFT graphs at $T \neq 0$ we refer the reader to the paper by Broglia *et al.* in ref. [24]). In the present work, eqs. (16) are used to calculate the GDR damping in hot even–even spherical nuclei.

3. The connection between the FT-QPNM equations and the strength function method at finite temperature. The choice of parameters and details of calculations

We employ the standard strength function method [52]) that has successfully been applied in studying the giant resonance damping in cold nuclei in the zero temperature QPNM [31]). This method has also been developed in different contexts of the response functions [12,15,21]) or the random matrix ensemble [26]). Thus, we define the strength function at finite temperature by analogy with the zero temperature case [31])

$$b(\Phi, \eta) = \sum_{\nu} |\Phi_{J\nu}|^2 \rho(\eta - \eta_\nu), \tag{20}$$

where $\Phi_{J\nu}$ is the amplitude of the excited state (14). If the excitation of the state (14) proceeds through the thermal one-phonon components of the wave function, we have

$$\Phi_{J\nu} = \sum_i R_i^\nu(J) \mathcal{M}_{Ji} \tag{21}$$

with $R_i^\nu(J)$ being the thermal one-phonon amplitudes defined in eq. (14). The density function $\rho(\eta - \eta_\nu)$ is parametrized as a lorentzian [31])

$$\rho(\eta - \eta_\nu) = \frac{1}{2\pi} \frac{\gamma_\nu}{(\eta - \eta_\nu)^2 + (\frac{1}{2}\gamma_\nu)^2}. \tag{22}$$

The matrix elements $\mathcal{M}_{Ji}$ are defined by the process under consideration. In our case they are the matrix elements of the E1 transitions from the thermal ground state (12) to the thermal one-phonon states [19]). In a form convenient to numerical calculations the strength functions (20) have been performed as [31,52])

$$b(\Phi, \eta) = \frac{1}{\pi} \mathrm{Im} \left\{ \sum_{ii'} \mathcal{A}_{ii'}(\eta + \tfrac{1}{2}i\gamma_\eta) \mathcal{M}_{Ji} \mathcal{M}_{Ji'} / \mathcal{F}(\eta + \tfrac{1}{2}i\gamma_\eta) \right\} \tag{23}$$

where $\mathcal{F}(\eta + \frac{1}{2}i\gamma_\eta)$ is the determinant in the l.h.s. of eq. (16) at complex values of energy and $\mathcal{A}_{ii'}$ is its minor. This determinant is calculated in the study of the giant resonance damping instead of solving the secular equations (16).

The centroid energy $\bar{E}$ and the spreading width $\Gamma^\downarrow$ of the strength function (23) are calculated in the energy interval $E_1 \leqslant \eta \leqslant E_2$ in a standard way [31])

$$\bar{E} = N^{-1} \int_{E_1}^{E_2} b(\Phi, \eta) \eta \, d\eta,$$

$$\Gamma^\downarrow = \left\{ 2.35 N^{-1} \int_{E_1}^{E_2} b(\Phi, \eta)(\eta - \bar{E})^2 \, d\eta \right\}^{1/2},$$

$$N = \int_{E_1}^{E_2} b(\Phi, \eta) \, d\eta. \tag{24}$$

In practice the energy interval $E_1 \leqslant \eta \leqslant E_2$ is chosen to be nearly the energy region where the experimental giant resonance is localized. Nevertheless, the comparison with the experimental width should be taken with some caution since the experimental data are derived as the lorentzian parameters, namely the full width at half maximum (FWHM) $\Gamma^\downarrow$ and its energy position $\bar{E}$. Having in mind the width $\Gamma^\downarrow$

from (24), we have chosen in the present work the interval (E_1, E_2) to be IV-GDR region for cold nuclei ^{58}Ni and ^{90}Zr, namely $8 \leq \eta \leq 24$ MeV for ^{58}Ni and $10 \leq \eta \leq 22$ MeV (see sect. 4) for ^{90}Zr. This energy interval is fixed at finite temperature. Due to eqs. (16) all the quantities included in eqs. (20)–(24) do depend now upon the temperature and the average values (13) are always understood.

We now focus on the physical insight into parameter γ_ν in eqs. (22)–(23) at finite temperature. It can be viewed as the damping of the thermal one-phonon state (13). Its inverse value $(\gamma_\nu)^{-1}$ is therefore understood as the hopping time needed for the transition from the thermal ground state to the thermal one-phonon excited state [cf. ref. [26])]. In the zero temperature QPNM the parameter γ_ν is state-independent $(\gamma_\nu \equiv \delta)$ and defines the way of representing the calculated results. It should be larger than an average distance between two-phonon states and much less than the region of location of a physical quantity calculated. In spherical cold nuclei $\gamma_\nu \equiv \delta$ has usually been chosen in the interval $0.1 \leq \delta \leq 1$ MeV. At finite temperature T, the parameter γ_ν can depend also upon T. In the theory of Fermi-liquids the damping of giant resonances is described as the absorption of the nuclear zero sounds. The absorption coefficient γ for a single GDR mode with energy η is found analytically from the Landau collisional integral as [53])

$$\gamma = a[\eta^2 + 4\pi^2 T^2] . \tag{25}$$

In eq. (25) the first term $\sim \eta^2$ in the r.h.s. corresponds to the damping of the GDR in cold systems whereas the second term contains a T^2-dependence of the damping at finite temperature. As the damping of the single-particle motion is not yet well understood we assume here that the parameter γ for every microscopic thermal phonon state with energy $\eta(T)$ has analogy with eq. (25). Moreover, as our calculations are performed only for the giant resonance region we put the energy dependence in eq. (25) to be on average $\eta \approx \eta^{\mathrm{GDR}}(T)$ where $\eta^{\mathrm{GDR}}(T)$ is the GDR energy at temperature T. Therefore, we have for γ the simple relation

$$\gamma_\eta \approx \delta[1 + (2\pi T/\eta^{\mathrm{GDR}}(T))^2] , \tag{26}$$

with $\delta \approx a[\eta^{\mathrm{GDR}}(T)]^2$. In our calculations at $T \neq 0$, eq. (26) should also be used. We note, however that the full microscopic knowledge of the parameter γ_η at $T \neq 0$ is still absent. In general, more detailed investigations are required. In the finite surface dominated systems, for example, a linear dependence on η and T for γ in eq. (25) has been suggested [24]). As is seen in refs. [13,19,42]), although there is a larger splitting for giant resonance modes in the FT-RPA, we still have a narrowing of the strength distribution with respect to the observed widths. The Landau splitting has not changed so much [19]) even with taking into account the temperature dependence of damping of γ-lines as in eq. (25) [ref. [42])]. This means the Landau damping alone is not enough to reproduce the experimental one, as has been mentioned in the Introduction. On the other hand, our approach in the FT-QPNM indeed can estimate the spreading width $\Gamma^\downarrow$ connected with quantal fluctuations. In this work we shall consider how a combination of the two mechanisms can influence the GDR damping in hot even–even spherical nuclei.

Let us now summarize briefly the choice of parameters we use in our approach and some details of calculations. At zero temperature the form of the mean-field in the QPNM is postulated by the Woods–Saxon potential W, whose parameters have been defined following ref. [54]). The single-particle energies calculated in this potential and employed in the QPNM calculations include discrete and quasidiscrete states arising from the centrifugal and Coulomb barriers. This truncation of the basis is apparently a good approximation for description of state characteristics up to energies ~25 – 30 MeV in medium and heavy nuclei. In any case the calculations performed in a full basis [55]) and with taking into account only bound and quasibound levels [56]) have given similar results for lowest ($2_1^+, 3_1^-, 4_1^+, 5_1^-$) states as well as for IV-E1 and IS-E2, E3 giant resonances. This has also been pointed out by Vdovin and Soloviev in ref. [57]). The radial part of the single-particle matrix elements $f_{jj'}^{(\lambda)}$ is described preferably by $R_\lambda(r) = \partial W(r)/\partial r$ [ref. [57])]. In principle, the phonon space in QPNM includes all multipolarities $\lambda \geq 1$. However, the number of collective phonons decreases drastically with increasing multipolarity λ. On the other hand, we do not take into account the continuum. This allows us to truncate the multipole phonon space to retain in our present calculations only the multipolarities up to $\lambda = 5$. For each multipolarity we take into account the whole one-phonon spectrum up to $\omega_\lambda \approx 24$ MeV. In constructions of two-phonon configurations in order to exclude the spurious states violating the Pauli principle we require at least one of two intermediate phonons be collective. It is the effective procedure for taking into account the Pauli principle approximately which has been mentioned above and discussed in detail in ref. [51]). Together with the approximation leading to eq. (16) at $T \neq 0$ discussed above this makes the very complicated computation task realizable, incorporating at the same time quantitatively physical effects of interest. The separable multipole forces are included with the multipole isoscalar $\varkappa_0^{(\lambda)}$ and isovector $\varkappa_1^{(\lambda)}$ constants [31]). The method of choosing these multipole constants has been presented previously and discussed thoroughly in many works within the QPNM framework [see e.g. refs. [31,57])]. A detailed review of the prescription for these parameters has been given recently by Galès, Stoyanov and Vdovin [58]). Its repetition here should be therefore superfluous. These authors have also pointed out that on average the spin–multipole interaction influences the fragmentation of single-particle states weakly. We do not therefore include the spin–multipole forces at all in our calculations. In general, at finite temperature, all the parameters of the potential, the coupling constants and, consequently, the single-particle spectrum etc., do depend upon the temperature. However, several numerical estimations [59]) have shown that the temperature dependence of the single-particle energies is rather smooth and weak up to $T \approx 6$ MeV. For this reason at the temperature of interest ($T < 6$ MeV) its effect can be neglected and we shall use the same single-particle energies E_j defined in the zero-temperature Woods–Saxon potential. Concerning

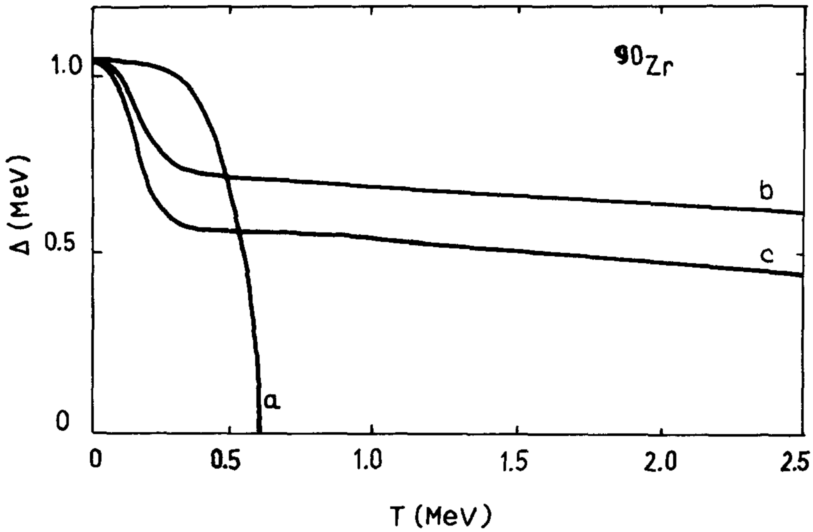

Fig. 2. Pairing gap versus temperature for ^{90}Zr: (a) FT-BCS gap Δ_T from eq. (2); (b) average gap $\langle\Delta_T\rangle$ from eq. (5) with the Δ-dependence of the mass $B(\Delta_T)$ omitted; (c) average gap $\langle\Delta_T(B)\rangle$ from eq. (5) with $B(\Delta_T)$ given in ref. [40]).

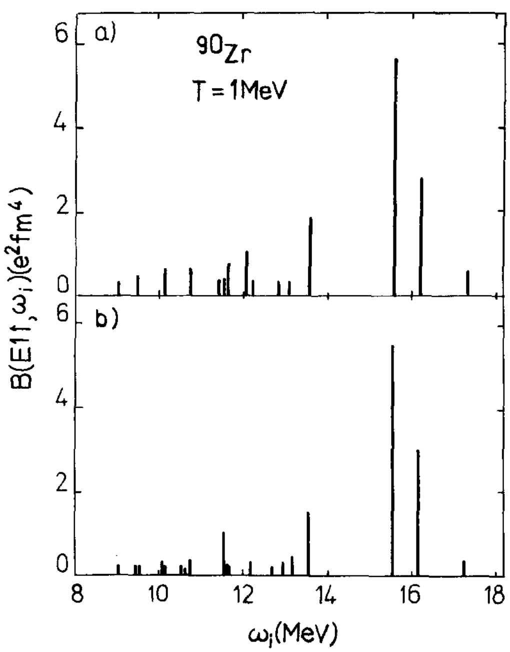

Fig. 3. Isovector dipole reduced transition probabilities $B(\mathrm{E}1\uparrow;\, \omega_{\lambda i})$ for ^{90}Zr calculated at temperature $T = 1$ MeV with (a) the gap $\langle\Delta_T\rangle$ (fig. 1b); (b) the gap $\langle\Delta_T(B)\rangle$ (fig. 1c).

the effective multipole interactions, the question about their temperature dependence is still open apart from some efforts to study the behaviour of the isoscalar multipole constants $\varkappa_0^{(\lambda)}$ as a functions of temperature [11,60]) that have shown nevertheless their very weak temperature dependence. We therefore shall also use at finite temperature the same values of the constants $\varkappa_{0,1}^{(\lambda)}$ defined at zero temperature by the above-mentioned method [cf. ref. [19])].

The thermal average superfluid pairing gaps computed from eqs. (5) and (6) in ^{90}Zr are plotted against temperatures in fig. 2. Although the Δ-dependence of the mass $B(\Delta_T)$ [ref. [40])] noticeably reduces the gap $\langle \Delta_T \rangle$ to the values $\langle \Delta_T(B) \rangle$, its effect on the strength distribution of the electric isovector dipole transitions computed from a formula in the FT-RPA framework [19]) is negligible (fig. 3). We therefore shall also reject this Δ-dependence of the mass everywhere throughout our calculations of the GDR dampings at finite temperature by omitting $B(\Delta_T)$ in eq. (6).

4. Numerical calculations

In the following, the strength functions $b(\mathrm{E1}, \eta)$ for the IV-GDR in even–even spherical nuclei ^{58}Ni and ^{90}Ar are calculated at several temperatures. For methodical consideration the calculations have been performed with several values of the parameter γ (eq. (22)). Namely, it has been taken temperature-independent $\gamma = \delta$ or temperature-dependent as in eq. (26). The effect of nonvanishing superfluid pairing gap at finite temperature on the strength distribution can be seen by comparing the results calculated by using the pairing gap (5) and with the temperature-independent pairing gap $\Delta_T = \Delta_{T=0}$.

The strength functions $b(\mathrm{E1}, \eta)$ are displayed in fig. 4 at $T = 0$, 1 and 3 MeV for ^{58}Ni. The results have been performed with the temperature-dependent gap (5) and with the temperature-independent $\gamma = \delta$ which takes the value 1 MeV in fig. 4a and 0.5 MeV in fig. 4b. The IV-GDR strength is distributed over a considerably wide energy interval from about 8 MeV to 24 MeV. The spreading width $\Gamma^{\downarrow}$ at zero temperature computed from eqs. (24) is 4.61 MeV. The IV-GDR centroid energy $\bar{\mathrm{E}}1$ is 17 MeV [cf. ref. [61])]. At finite temperature the centroid energy $\bar{\mathrm{E}}1$ is shifted to lower values while the width $\Gamma^{\downarrow}$ increases slightly. At $T = 3$ MeV the value of $\Gamma^{\downarrow}$ is 5.7 MeV while the centroid energy $\bar{\mathrm{E}}1$ is localized at 15.2 MeV. The energy-weighted sum rule (EWSR) for the IV GDR in our calculations is exhausted by 85–90% and is considerably stable with varying temperature. By comparing fig. 4a and fig. 4b we can see how a larger value of $\gamma = \delta$ smooths the finite structure of the IV-GDR. The calculations using a smaller $\delta = 0.5$ MeV also indicate a clear enhancement of the lower lying GDR modes at finite temperature (fig. 4b). As concerns the role of the superfluid pairing interaction, we find it to effect very slightly the values of the IV-GDR spreading width $\Gamma^{\downarrow}$ as well as the centroid energy $\bar{\mathrm{E}}1$. In fact, the calculations performed with the gap $\Delta_T = \Delta_{T=0} \approx 1.4$ MeV in ^{58}Ni give even at $T = 3$ MeV the value 4.58 MeV for $\Gamma^{\downarrow}$ and 15.6 MeV for $\bar{\mathrm{E}}1$ of the

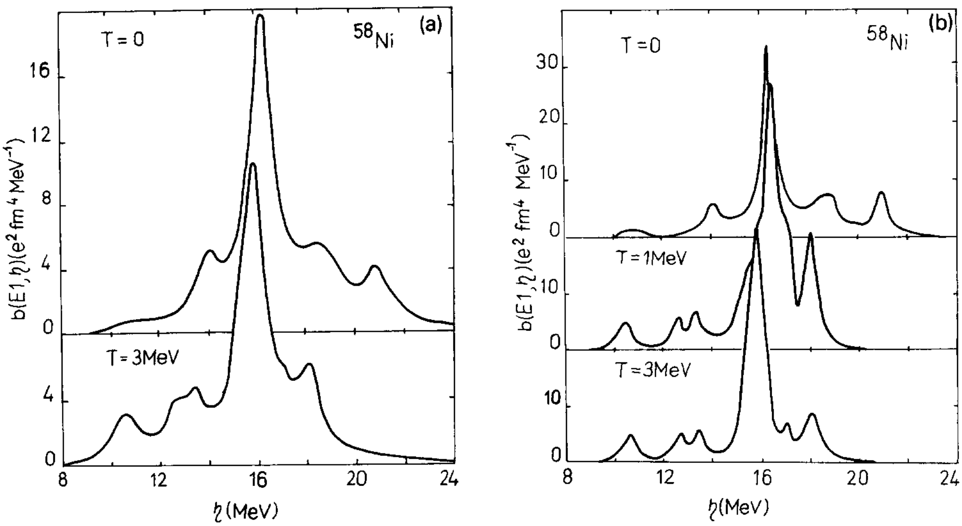

Fig. 4. IV-GDR strength functions $b(\mathrm{E1}, \eta)$ at several temperatures for ^{58}Ni calculated with the gap $\langle \Delta_T \rangle$. Results are displayed with the parameter $\gamma \equiv \delta$: (a) $\delta = 1$ MeV; (b) $\delta = 0.5$ MeV.

IV-GDR. However, the forms of the strength distribution in the two cases (with the temperature dependent average gap $\langle \Delta_T \rangle$ and with $\Delta_T = \Delta_{T=0}$) are somewhat different as can be clearly seen by comparing fig. 4 and fig. 5. The strengths in the case with $\Delta_T = \Delta_{T=0}$ are more concentrated around the centroid energy (fig. 5) and the shift this centroid energy with increasing T is rather weak. In general, the calculated fragmentation of the IV-GDR does not change much with varying temperature.

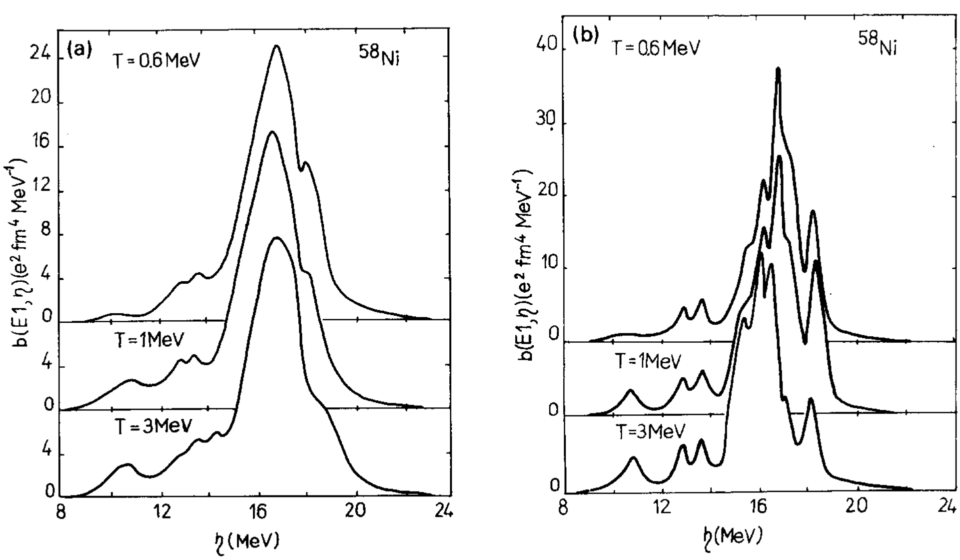

Fig. 5. The same as in fig. 4 with the temperature independent gap $\Delta_T \equiv \Delta_{T=0}$.

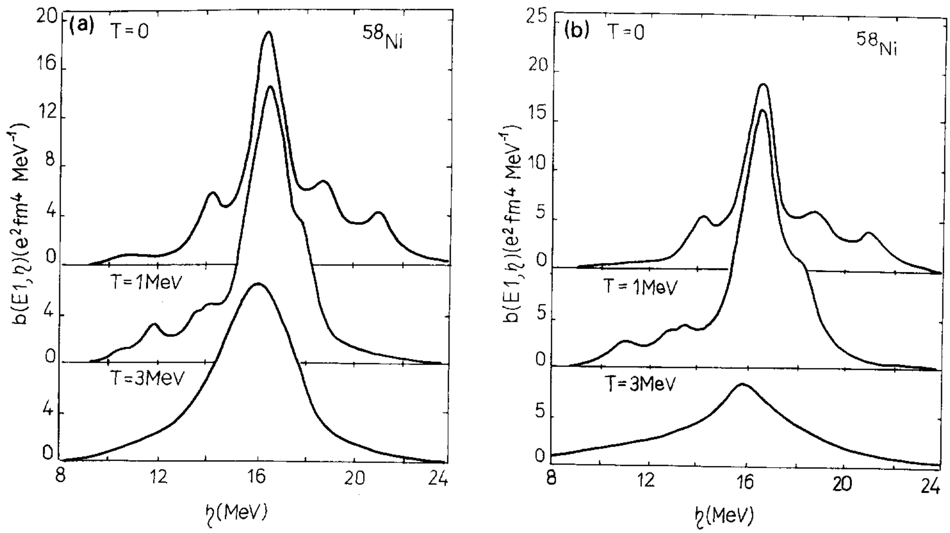

Fig. 6. The same as in fig. 4 with γ given by eq. (26) and (a) $\Delta = \Delta_{T=0}$; (b) $\Delta = \langle \Delta_T \rangle$.

By taking into account the temperature-dependent parameter $\gamma(T)$ from eq. (26) we find its value to increase with temperature T. Therefore, the Landau damping of this kind tends to smooth out the finite structures of the GDR while leaving the strong collective states visible (fig. 6). At $T = 3$ MeV the effect of the temperature dependence of $\gamma(T)$ is so strong that it smears out all the finite structures in a single large bump localized in the region of the IV-GDR. The form of the strength distribution in the considered energy interval changes noticeably. Nevertheless, for $T = 3$ MeV the spreading width $\Gamma^{\downarrow}$ in this energy interval increases weakly as compared to the case with the temperature independent $\gamma = \delta$.

The widths $\Gamma^{\downarrow}$ and the centroid energies $\bar{\mathrm{E}}1$ are plotted against temperature T in fig. 7 for ^{58}Ni. The results displayed in this figure clearly show the effect of the nonvanishing pairing gap at finite temperature on these giant resonance characteristics. In fact, at each value of T, the width $\Gamma^{\downarrow}$ calculated with $\Delta_T = \Delta_{T=0}$ is always smaller than the one calculated with the gap $\langle \Delta_T \rangle$ whereas the corresponding centroid energy $\bar{\mathrm{E}}1$ is always higher. The tendency of the centroid energy $\bar{\mathrm{E}}1$ to shift down and the broadening of $\Gamma^{\downarrow}$ observed in experiments with increasing T are also evident. However, the broadening turns out to be rather small at finite T and systematically much lower than the experimental values obtained in this mass number region with the same given value of T (excitation energy E^*) [7]. This situation is not surprising. Indeed our calculations have been performed for spherical hot nuclei without the angular momentum effects, while the experimental data have been obtained in highly excited nuclei with a given temperature T and a given grazing angular momentum l_0 [refs. [6-10]]. Therefore, we can conclude that the contribution of the interaction

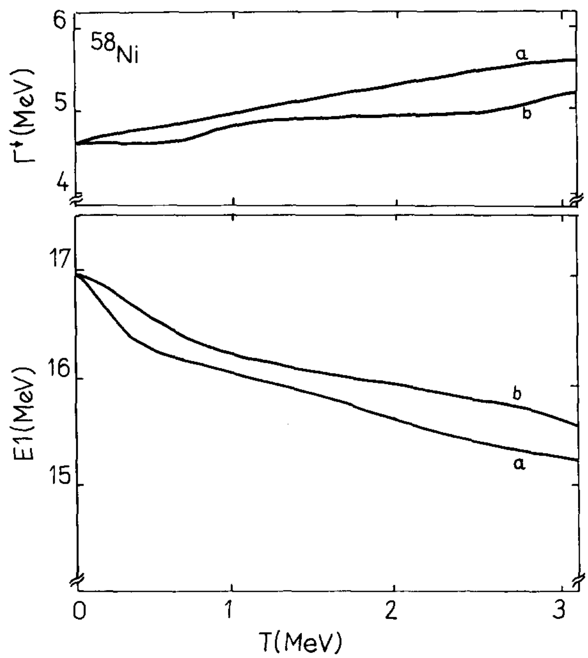

Fig. 7. IV-GDR spreading width $\Gamma^{\downarrow}$ and centroid energy $\bar{\mathrm{E}}1$ for ^{58}Ni versus temperature. Results are displayed for and (a) $\Delta = \langle\Delta_T\rangle$; (b) $\Delta = \Delta_{T=0}$.

with (2p2h) configurations to the sizeable spread of the GDR at finite temperature changes weakly, while the shape deformation leading to the presence of spin could be the principal cause.

A similar situation is seen in ^{90}Zr nucleus. Fig. 8 represents the strength functions calculated at several temperatures with the temperature dependent averaging pairing gap $\langle\Delta_T\rangle$ in ^{90}Zr for two cases: with the temperature independent parameter $\gamma = \delta =$ 1 MeV with $\gamma(T)$ given by eq. (26). The IV-GDR characteristics in this nucleus are more stable with varying temperature T, as compared to the case of the hot ^{58}Ni. The IV-GDR spreading widths $\Gamma^{\downarrow}$ and the centroid energies $\bar{\mathrm{E}}1$ in ^{90}Zr at $T = 0, 1, 2$ and 3 MeV are collected in table 1. At zero temperature $\Gamma^{\downarrow}$ is 3.7 MeV and $\bar{\mathrm{E}}1$ is 16.23 MeV [cf. ref. [62])]. At finite temperature, while the centroid energy $\bar{\mathrm{E}}1$ is shifted down noticeably, the spreading with $\Gamma^{\downarrow}$ increases very smoothly and weakly with increasing temperature T. The experimental data have been obtained in the highly excited ^{90}Zr [ref. [8])], where a strongly broadened FWHM $\Gamma \approx 8.8$ MeV has been observed at $T \approx 1.7$ MeV and grazing angular momentum $l_0 \approx 5\hbar$. The discrepancy between them and our theoretical microscopic calculations taking into account the couplings to (2p2h) states at finite temperature without rotation demonstrates again the important role of the shape deformation in the broadening of GDR in highly excited nuclei.

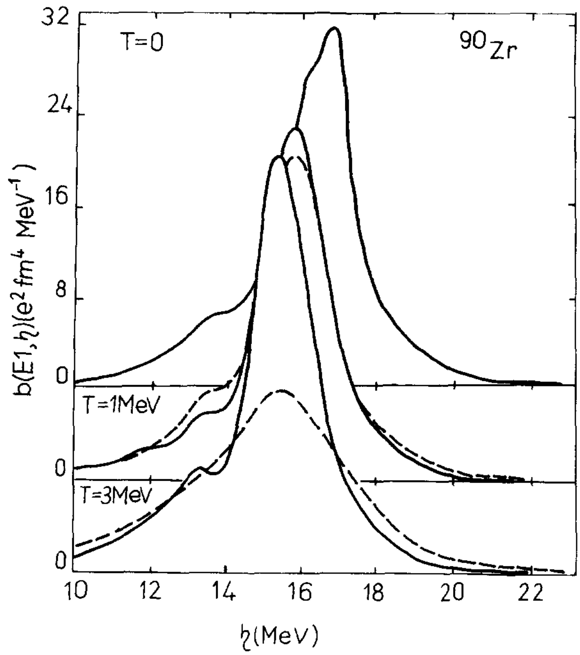

Fig. 8. The same as in fig. 4 for ^{90}Zr. Solid curves denote the results calculated with $\gamma = 1$ MeV while the dashed curves correspond to the results obtained with γ given by eq. (26).

Qualitatively, our numerical calculations give a similar conclusion as compared to the calculations performed in the finite temperature Matsubara formalism taking into account the couplings "1p1h$\oplus$phonon" in the NFT [24]). This indicates the main contribution of the one-phonon collective degrees of freedom in the damping of giant resonances at zero as well as at finite temperature in the case with the rotation omitted. On the other hand, our approach neatly gives a larger increase in the IV-GDR spreading width with T. This fact of course is germane to the contribution of different graphs in both the approaches. Indeed, as has been dis-

TABLE 1

IV-GDR spreading widths $\Gamma^{\downarrow}$ and centroid energies $\bar{\mathrm{E}}1$ calculated with the average pairing gaps $\langle \Delta_T \rangle$ and $\gamma = 1$ MeV at several temperatures for ^{90}Zr. (All the values are given in MeV)

T	$\Gamma^{\downarrow}$	$\bar{\mathrm{E}}1$
0	3.70	16.23
1	4.11	15.53
2	4.42	15.20
3	4.60	15.01

cussed above, our calculations include the graph in fig. 1a, which contains as a special case the most important NFT graphs in fig. 1d. However, we have neglected 1e, f in the finite temperature Matsubara formalism of the NFT [24]). Although the contribution of the graphs 1b, c is assumed to be small, their incorporation would alter slightly the results. A similar broadening of the IV-GDR at moderate temperature has also been obtained recently within the framework of a semiclassical approach based on the Landau–Vlasov equations suggested in ref. [27]) for hot nuclei without rotation. This comparison can shed light on the question which still remains to be answered about the relation between the microscopic approaches including those of the present work and ref. [24]), and the semiclassical ones. Of course, we have neglected in our calculations the temperature dependence of the interactions, the effect of phonon ground-state correlations and phonon scattering at finite temperature etc., as has been mentioned above. However, the investigations in hot and strongly rotating nuclei [23,25,29,30]) together with the experimental data [1-10]) obviously confirm the key role of the angular momentum effects in the description of the broadening of the IV-GDR width in highly excited nuclei.

5. Conclusions

In this work to study the damping of the IV-GDR in hot even–even spherical nuclei, we have applied our approach, the FT-QPNM, which extends the QPNM to finite temperature. This model therefore allows us to take consequently the couplings to (2p2h) configurations into account at finite temperature. Summarizing the numerical results we find:

(i) Although the account of the couplings to (2p2h) states can reproduce rather well the damping width and the centroid energy of IV-GDR at zero temperature, this is not sufficient to describe the large broadening of the IV-GDR damping width with increasing temperature T. The shift down of the IV-GDR centroid energy at finite T seems to be in reasonable agreement with the experimental data, while the spreading width $\Gamma^{\downarrow}$ (eq. (24)) in the considered energy interval increases not so much with T.

(ii) The zero-sound Landau damping effect, included in our calculations at the temperature dependence of the parameter γ from eq. (26), tends to smooth out the fine structure of the IV-GDR and leaves only a broad peak of the Lorentz form with increasing temperature.

(iii) The account of the nonvanishing average superfluid pairing gap due to thermal fluctuations enlarges slightly the damping width and reduces the centroid energy of IV-GDR at each value of the temperature.

Therefore, we conclude that the contribution of the couplings to (2p2h) configurations to the IV-GDR spreading width does not change much with temperature $T \leqslant 5$ MeV. To describe the data (i.e. the sizeable width of the IV-GDR at $T \neq 0$) one has to take into account at least:

(i) deformations together with the angular momentum effect.

(ii) the effects of phonon ground-state correlations, phonon scattering and the continuum which can be noticeable at moderate temperatures;

(iii) the couplings to states more complicated than 2p2h.

The temperature-dependence of residual interactions is also a problem that remains to be studied in detail.

It is worth recognizing that the width and structure of the GDR at high temperatures ($T > 5$–6 MeV) may not follow the pattern observed at lower T. Indeed a narrow GDR has been implied from the data [9]), suggesting a different coupling mechanism of the GDR to deformation. Recently, a simple model of motional narrowing has been suggested by Broglia *et al.* [26]). As this phenomenon takes place in the high-temperature region and cannot be considered without deformation, it goes beyond the framework of our present study. A microscopic quantitative investigation of this effect is of course highly desirable.

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References

1) J. Barrette and A.M. Sandorfi, Comm. Nucl. Part. Phys. XII (1983) 57

2) J.O. Newton, B. Herskind, R.M. Diamond, E.L. Dines, J.E. Draper, K.H. Lindenberger, C. Schuck, S. Shih and F. Stephens, Phys. Rev. Lett. **46** (1981) 1383

3) W. Hennerici, V. Metag, H.J. Heinrich, R. Repnov, W. Wahh, D. Habs, K. Helmer, U.V. Helmolt, H.W. Heyng, B. Kolb, D. Pelte, D. Schwalm, R.S. Simon and R. Albrech, Nucl. Phys. **A396** (1983) 329c

4) E.F. Garman, K.A. Snover, S.H. Chao, S.K.B. Hesmondhalgh, W.N. Catford and P.M. Walker, Phys. Rev. **C28** (1983) 2554

5) J.E. Draper, J.O. Newton, L.G. Sobotka, K.H. Lindenberger, G.J. Wozniak, L.G. Moretto, F.S. Stephens, R.M. Diamond and R.J. McDonald, Phys. Rev. Lett. **49** (1982) 434

6) C.A. Gossett, K.A. Snover, J.A. Behnr, G. Feldman and J.L. Osborne, Phys. Rev. Lett. **54** (1985) 1486; J.A. Behr, G. Feldman, H. Glatzel, C.A. Gossett, J.H. Gundlach, M. Kicinska-Habior and K.A. Snover, Nuclear Physics Laboratory Annual Report, University of Washington (1986) p. 10; (1987) p. 10

7) J.A. Behr, G. Feldman, C.A. Gossett, J.H. Gundlach, M. Kicinska-Habior and K.A. Snover, Nuclear Physics Laboratory Annual Report, University of Washington (1986) p. 4

8) J.A. Behr, G. Feldman, J.H. Gundlach, C.A. Gossett and K.A. Snover, *ibid* (1986) p. 12;
J.A. Behr, G. Feldman, C.A. Gossett, J.H. Gundlach, M. Kicinska-Habior and K.A. Snover, *ibid* (1987) p. 12; Contributed paper in First Topical Meeting on giant resonance excitation in heavy-ion collisions, Legnaro (Padova) Italy, Sept. 21-25, 1987
9) J.J. Gaardhøje, O. Anderson, R.M. Diamond, C. Ellegaard, L. Grodzins, B. Herskind, Z. Sujkowski and P.M. Walker, Phys. Lett. **B139** (1984) 273;
J.J. Gaardhøje, C. Ellegaard, B. Herskind and S.G. Steadman, Phys. Rev. Lett. **53** (1984) 148;
J.J. Gaardhøje, A.M. Bruce and B. Herskind, Nucl. Phys. **A482** (1988) 121c;
D.R. Chakrabarty, S. Sen, N. Alamanos, P. Paul, R. Shicker, J. Stachel, M. Thoennessen and J.J. Gaardhøje, Phys. Rev. **C36** (1987) 1886;
D.R. Chakrabarty, Nucl. Phys. **A382** (1988) 81c
10) K.A. Snover, Ann. Rev. Nucl. Part. Sci. **36** (1986) 545
K.A. Snover, Nucl. Phys. **A482** (1988) 13c
11) A.V. Ignatyuk, Izv. Acad. Nauk. SSSR (ser. fiz.), **38** (1974) 2613;
M.I. Baznat, A.V. Ignatyuk and N.I. Pyatov, Yad. Fiz. **30** (1979) 949;
A.V. Ignatyuk, Statistical properties of excited atomic nuclei (Energoatomizdat, Moscow, 1983) (in Russian)
12) M. Sommermann, Ann of Phys. **151** (1983) 163;
K. Tanabe, Phys. Rev. **C37** (1988) 2802
13) D. Vautherin and N. Vinh Mau, Phys. Lett. **B120** (1983) 261; Nucl. Phys. **A422** (1984) 140
14) J. Meyer, P. Quentin and B. Brack, Phys. Lett. **B133** (1983) 279
15) O. Civitarese, R.A. Broglia and C.H. Dasso, Ann. of Phys. **156** (1984) 142
16) J.L. Egido and H.A. Weidenmüller, Phys. Lett. **B208** (1988) 58;
M. Barranco, A. Polls and J. Martorrel, Nucl. Phys. **A444** (1985) 445;
A. Milian, M. Barranco, D. Mas and R.J. Lombard, Z. Phys. **A330** (1988) 5
17) H. Sagawa and G.F. Bertsch, Phys. Lett. **B146** (1984) 138
J. Bar-Touv, Phys. Rev. **C32** (1985) 1369
18) A. Kerman and S. Levit, Phys. Rev. **C24** (1981) 1029;
A. Kerman, S. Levit and T. Troudet, Ann of Phys. **154** (1984) 456
19) Nguyen Dinh Dang, J. of Phys. **G11** (1985) L125
20) F. Alasia, O. Civitarese and M. Reboiro, Phys. Rev. **C35** (1987) 812
21) G.F. Bertsch, P.F. Bortignon, R.A. Broglia and C.H. Dasso, Phys. Lett. **B80** (1979) 161;
P.F. Bortignon and R.A. Broglia, Nucl. Phys. **A371** (1981) 405;
G.F. Bertsch, P.F. Bortignon and R.A. Broglia, Rev. Mod. Phys. **55** (1983) 287
22) V.V. Voronov, Nguyen Dinh Dang, V.Yu. Ponomarev, V.G. Soloviev and Ch. Stoyanov, Yad. Fiz. (Sov. J. Nucl. Phys) **40** (1984) 683
23) M. Gallardo, M. Diebel, T. Døssing and R.A. Broglia, Nucl. Phys. **A443** (1985) 415
24) P.F. Bortignon, R.A. Broglia, G.F. Bertsch and J. Pacheco, Nucl. Phys. **A460** (1986) 149;
R.A. Broglia, D.R. Bes, P.F. Bortignon and M. Gallardo, Nuclear structure, reactions and symmetries; 5-14 June (1986), Dubrovnik, ed. R.A. Meyer and V. Paar, vol. 2 (World Scientific, Singapore, 1986) p. 882
25) M. Gallardo, F.J. Luis and R.A. Broglia, Phys. Lett. **B191** (1987) 332
26) R.A. Broglia, W.E. Ormand and M. Borromeo, Nucl. Phys. **A482** (1988) 141c
27) Cai Yanhuang and M. Di Toro Phys. Rev. **C39** (1988) 105
28) P. Ring, L.M. Robledo, J.L. Egido and M. Faber, Nucl. Phys. **A419** (1984) 261;
P. Ring, Nucl. Phys. **A482** (1988) 27c
29) A.V. Ignatyuk, I.N. Mikhailov, L.H. Molina, R.G. Nazmitdinov and K. Pomorsky, Nucl. Phys. **A346** (1980) 191;
S.N. Fedotkin, I.N. Mikhailov and R.G. Nazmitdinov, Phys. Lett. **B121** (1983) 15;
J. Kvasil and R.G. Nazmitdinov, Acta Universitatis Carolinae Mathematica et Physica **29** (1988) 33
30) B. Lauritzen, T. Døssing and R.A. Broglia, Nucl. Phys. **A457** (1986) 61;
R.A. Broglia, T. Døssing, B. Lauritzen and B.R. Mottelson, Phys. Rev. Lett. **58** (1987) 326
31) V.G. Soloviev, Theory of complex nuclei (Pergamon, Oxford, 1976);
V.G. Soloviev, Fiz. Elem. Chastits At. Yadra (Sov. J. Part. Nucl.) **23** (1978) 580;
V.G. Soloviev, Prog. Part. Nucl. Phys. **19** (1987) 107

32) Nguyen Dinh Dang, Teor. Mat. Fiz. (Sov. J. Theor. Mat. Phys.) **64** (1985) 69
33) Nguyen Dinh Dang, Phys. Math. Candidate Dissertation (Sov. Ph. D. thesis), JINR, 4-85-38 (UDK 539.142/143) (Dubna 1985) (in Russian)
34) V.V. Voronov and Nguyen Dinh Dang, Izv. Acad. Nauk, SSSR (ser. fiz.) **48** (1984) 857
35) Nguyen Dinh Dang, Z. Phys. **A327** (1987) 41
36) C. Yannouleas and S. Jang, Nucl. Phys. **A455** (1986) 40
37) C. Yannouleas, M. Dworzecka and J.J. Griffin, Nucl. Phys. **A379** (1982) 256; **A397** (1983) 239
38) M. Abe and N. Takigawa, Phys. Lett. **B209** (1988) 149
39) L.G. Moretto, Nucl. Phys. **A185** (1972) 145; **A226** (1974) 9; Phys. Lett. **B40** (1972) 1; **B44** (1973) 494; A.L. Goodman, Phys. Rev. **C29** (1984) 1887
40) J.L. Egido, P. Ring, S. Iwasaki and H.J. Mang, Phys. Lett. **B154** (1985) 1
41) Nguyen Dinh Dang, JINR Communication, E4-88-40, (Dubna, 1988)
42) Nguyen Dinh Dang and Nguyen Zuy Thang, J. of Phys. **G14** (1988) 1471
43) R. Rossignoli, A. Plastino and J.P. Vary, Phys. Rev. **C37** (1988) 314
44) Nguyen Dinh Dang and Nguyen Zuy Thang, JINR preprint P4-88-660 (Dubna, 1988)
45) Nguyen Dinh Dang and V.Yu. Ponomarev, Izv. Acad. Nauk SSSR (ser. fiz.), **47** (1983) 2147
46) L.D. Landau and E.M. Lifshitz, Course of theoretical physics, vol. 5: Statistical physics (Nauka, Moscow, 1964) (in Russian)
47) M.K. Volkov, A. Pavlikovsky, V. Rybarska and V.G. Soloviev, Izv. Acad. Nauk. SSSR (ser. fiz). **27** (1963) 878;
I.N. Mikhailov, JETP (Sov. J. Exp. Theor. Phys). **45** (1963) 1102
48) Y. Nogami, Phys. Rev. **134** (1964) B313; Phys. Lett. **15** (1965) 335;
Y. Nogami and I.J. Zucker, Nucl. Phys. **60** (1964) 203;
H.C. Pradhan, Y. Nogami and J. Law, Nucl. Phys. **A201** (1973) 357
49) Nguyen Dinh Dang, JINR preprint, E4-88-307 (Dubna, 1988)
50) Nguyen Dinh Dang and V.Yu. Ponomarev, Teor. Mat. Fiz. (Sov. J. Theor. Math. Phys.) **57** (1983) 154
51) V.Yu. Ponomarev, Ch. Stoyanov and O. Stoyanova, JINR Communication, P4-81-704 (Dubna, 1981)
52) A. Bohr and B.R. Mottelson, Nuclear structure, vol. I (North-Holland, Amsterdam, 1969)
53) D. Pines and P. Nozieres, The theory of quantum liquids, vol. I: Normal Fermi liquids (Benjamin, New York, 1966);
L.D. Landau and E.M. Lifshitz, Theoretical physics, vol. 10: Physical cinetics (E.M. Lifshitz and L.P. Pytaevsky) (Nauka, Moscow, 1979) p. 383
54) V.A. Chepurnov, Yad. Fiz. (Sov. J. Nucl. Phys) **6** (1967) 955;
K. Takeuchi and P.A. Moldauer, Phys. Lett. **B28** (1969) 384;
V.G. Soloviev, Ch. Stoyanov and A.I. Vdovin, Nucl. Phys. **A288** (1977) 376
55) V.V. Palchik, N.I. Pyatov and S.A. Fayans, Yad. Fiz. (Sov. J. Nucl. Phys.) **34** (1981) 903
56) V.V. Voronov and C.-Z. Khuong, Izv. Acad. Nauk SSSR (ser. fiz.) **45** (1981) 1910
57) A.I. Vdovin and V.G. Soloviev, Fiz. Elem. Chastists. At. Yadra (Sov. J. Part. Nuco.) **14** (1983) 237
58) S. Galès, Ch. Stoyanov and A.I. Vdovin, Phys. Reports **166** (1988) 125
59) M. Brack and P. Quentin, Phys. Lett. **B42** (1974) 159; Nucl. Phys. **A361** (1981) 35;
P. Bonche, S. Levit and D. Vautherin, Nucl. Phys. **A427** (1984) **178**
60) A.I. Blokhin and A.V. Ignatyuk, Problems of nuclear physics and cosmic rays (Vyshaya shcola, Kharkov, 1977) vol. 7, p. 111 (in Russian);
A.V. Ignatyuk and I.N. Mikhailov, Yad. Fiz. (Sov. J. Nucl. Phys.) **30** (1979) 665
61) D.C. Kocher and R.L. Auble, Nucl. Data Sheets **19** (1976) 445
62) B.L. Berman, J.T. Caldwell, R.R. Harvey, M.A. Kelly, R.L. Bramblett and S.C. Fultz, Phys. Rev. **162** (1967) 1098;
A. Leprêtre, H. Beil, R. Bergere, P. Carlos, A. Veyssière and M. Sugawara, Nucl. Phys. **A175** (1971) 609