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TEMPERATURE DEPENDENCE OF GIANT DIPOLE RESONANCE WIDTH

A. I. Vdovin, A. N. Storozhenko

Bogoliubov Laboratory of Theoretical Physics, Joint Institute for Nuclear Research, Dubna, Moscow region, Russia

The quasiparticle-phonon nuclear model extended to finite temperature within the framework of the thermo field dynamics is applied to calculate a temperature dependence of the spreading width $\Gamma^{\downarrow}$ of a giant dipole resonance. Numerical calculations are made for ^{120}Sn and ^{208}Pb nuclei. It is found that the width $\Gamma^{\downarrow}$ increases with T. The reason of this effect is discussed as well as a relation of the present approach to other ones existing in the literature.

Introduction

GDR was found in a hot rotating nucleus formed in a collision of two heavy ions as early as 1981 [1]. As a result of quite sophisticated experiments performed during 20 years some integral characteristics of GDR were carefully studied. In particular, it is well proved that the energy of GDR and the exhaustion of the model independent Energy Weighted Sum Rule are quite stable against temperature increase. At the same time one observes a strongly increasing width of GDR with temperature of a nucleus.

Several processes contribute to the GDR width at finite temperature [2 - 4]. Among them are quantum fluctuations which exist already in a cold nucleus: the Landau damping, the coupling with surface vibrations, the collisional damping (i.e. the coupling to incoherent two-particle – two-hole excitations) and the coupling to the single-particle continuum. At $T \neq 0$ the thermal fluctuations of a nuclear shape appear. Moreover, since a hot compound nucleus usually carries a large angular momentum, the rotation also affects the GDR width.

Extracting the GDR characteristics from the measured γ-spectra is not an absolutely unambiguous procedure. These spectra are in fact a weighted sum of the γ-ray yields emitted by many nuclei populated in the decay of the initial compound nucleus. The extracted GDR characteristics depend to some extent on assumptions about a shape of E1 strength function, and mass- and temperature-dependence of its parameters [5]. Also, the temperatures inferred from experimental excitation energy of a hot compound nucleus are sensitive to the level density parameter which is not known very accurately. In this respect, the impressive example is the fate of a phenomenon of the so-called saturation of the GDR width at $T \geq 3,5$ - 4 MeV. After the appearance of new data and reanalysis of the previous ones [6, 7] the GDR width Γ_{exp} was found permanently increasing up to $T \approx 3,2$ MeV. It was also established that the information about GDR at higher temperatures cannot be extracted reliably from the existing data.

Even a more ambiguous problem is the disentangling of different contributions to the experimental GDR width. Fortunately, due to the experiments with inelastically scattered α-particles which yield a compound system with a small angular momentum [8] the effects of rotation and temperature on the GDR width was separated. However, in most cases conclusions can be made only by comparing the final results of theoretical calculations with the measured experimental value. Sometimes conclusions appear to be controversial. For example, the adiabatic coupling model [9] reasonably describes the experimental data on the GDR width in ^{120}Sn and ^{208}Pb supposing the intrinsic GDR width Γ almost independent of temperature. According to studies [9], the main effect, which explains increasing of the total Γ_{exp}, is the thermal nuclear shape fluctuations. On the other hand, according to [10], the behavior of the GDR parameters in the compound nucleus ^{86}Mo cannot be explained by assuming the intrinsic width be a constant. Moreover, the recent measurement of the GDR width in ^{120}Sn at temperature $T \approx 1$ MeV [11] reveals an over-estimation of Γ_{exp} by the adiabatic coupling model.

Different theoretical approaches also predict a quite different T-dependence for the spreading GDR width $\Gamma^{\downarrow}$. The first calculations of the thermal behavior of $\Gamma^{\downarrow}$ were performed in [12]. At that time, it was already well known that the coupling of a single-particle motion with collective surface vibrations is the main mechanism of damping of giant resonances in cold nuclei. In [12], a temperature dependence of this coupling was studied with the Matsubara thermal Green's function technique and it was found that the GDR width was nearly constant when T increased. The physical ground of these calculations was the Nuclear Field Theory [13] treating a nucleus as a system of interacting quasiparticles and vibrations (RPA phonons). In more recent studies [14] the very weak dependence $\Gamma^{\downarrow}$ on T was explained by the cancellation effect between self-energy and vertex contributions. However, several years ago in [15], where the problem was studied within the same formalism and under the same physical assumptions as in [12, 14], an increment of the spreading GDR width with T was found.

The latter result is in correspondence with that of the approaches taking into account the coupling with incoherent 2p-2h excitations (the collisional damping) [4, 16, 17]. In most cases, calculations predict the increase of the GDR width with increasing in temperature; although the calculated width exhibits weaker temperature dependence and numerical predictions are quite sensitive to the effective nucleon-nucleon interaction used.

Thus, the current situation with the temperature dependence of the GDR spreading width is not clear. That is why new approaches to the problem are desirable. We present here results of the approach developed in [18 - 20] and based on the two main ingredients: the quasiparticle-phonon nuclear model (QPM) [21 - 23] and the thermo field dynamics (TFD) [24, 25]. The physical basis of QPM is very similar to that of the nuclear field theory, and both the models have produced quite close results as applied to nuclear structure calculations at T=0. In Refs. [18 - 20], the QPM was extended to finite temperatures by the use of the TFD formalism. Here we report the results of numerical calculations within the TFD-QPM approach. More detailed discussion of the subject can be found in [26].

QPM at finite temperature

Thermal RPA

We start with the QPM Hamiltonian which consists of phenomenological mean fields for protons and neutrons and separable multipole particle-hole interactions with both the isoscalar and isovector terms

$$H = \sum_{jm\tau}(E_j - \lambda_\tau)c^+_{jm}c_{jm} - \frac{1}{2}\sum_{\lambda\mu}\sum_{\tau,\rho=\pm1}(\chi_0^{(\lambda)} + \rho\chi_1^{(\lambda)})M^+_{\lambda\mu}(\tau)M_{\lambda\mu}(\rho\tau), \tag{1}$$

where $M^+_{\lambda\mu}(\tau)$ is the single-particle multipole operator

$$M^+_{\lambda\mu}(\tau) = \sum_{j_1m_1j_2m_2}\left\langle j_1m_1\middle|M_{\lambda\mu}\middle|j_2m_2\right\rangle c^+_{j_1m_1}c_{j_2m_2}.$$

Here c^+_{jm}, c_{jm} are creation and annihilation operators of particles and holes with quantum numbers n, l, j. m ≡ j, m. The index τ is an isotopic one and takes two values τ = n, p. The symbol Σ^τ means that the summation is taken only over neutron or proton single-particle (hole) states and changing the sign of τ means changing n↔p. The parameters $\chi_0^{(\lambda)}, \chi_1^{(\lambda)}$ are the coupling constants of the isoscalar and isovector multipole-multipole interactions, respectively.

To avoid unnecessary complications in the forthcoming formulae, we omit the pairing interaction in the model Hamiltonian (1)[1].Moreover, since we treat the pairing interaction in the thermal BCS approach the superfluid gap vanishes at T ≈ 1 MeV. Thus, the pairing correlations cannot play essential role in the temperature range where $\Gamma_{\exp}$ is measured.

The first step in treating nuclear dynamics governed by the model Hamiltonian (1) at finite temperatures is formal doubling of the Hilbert space of a nucleus. To this aim a fictitious (tilde-) nucleus which is of exactly the same structure as the initial one is introduced. Thus, the tilde creation and annihilation operators $\widetilde{c}^{+}_{jm}, \widetilde{c}_{jm}$ appear in the game. An excitation spectrum of a hot many-body system is obtained by diagonalization of the thermal Hamiltonian $\boldsymbol{H} = H - \widetilde{H}$. The thermal behavior of the system is controlled by the thermal vacuum state, which is the $\Gamma_{\exp}$ eigenstate of $\boldsymbol{H}$ with the zero eigenvalue.

Our starting point is the "mean field + RPA" scheme. To construct the Fock space of a hot nucleus, we make a unitary thermal Bogoliubov transformation from our initial and tilde operators c^{+}_{jm}, c_{jm}, $\widetilde{c}^{+}_{jm}, \widetilde{c}_{jm}$ to thermal quasiparticle operators $\beta^{+}_{jm}, \beta_{jm}, \tilde{\beta}^{+}_{jm}, \tilde{\beta}_{jm}$. The transformation has the form

$$\begin{aligned} \beta_{jm} &= \sqrt{1-n_j}\,c_{jm} - \sqrt{n_j}\,\tilde{c}^{+}_{jm} \\ \tilde{\beta}_{jm} &= \sqrt{1-n_j}\,\tilde{c}_{jm} + \sqrt{n_j}\,c^{+}_{jm} \end{aligned}, \qquad (2)$$

where n_j is a thermal Fermi occupation number

$$n_j = \frac{1}{1+\exp\left[(\mathrm{E_j}-\lambda_\tau)/\mathrm{T}\right]}.$$

The chemical potentials $\lambda_{n,p}$ are adjusted to fulfill demands of the neutron and proton number conservation in average. The transformation (2) saves the diagonal form of the single-particle part of the thermal Hamiltonian, i.e.

$$\boldsymbol{H}_{\mathrm{TSP}} = H_{\mathrm{SP}} - \tilde{H}_{\mathrm{SP}} = \sum_{jm\tau}\left(E_j - \lambda_\tau\right)\left(\beta^{+}_{jm}\beta_{jm} - \tilde{\beta}^{+}_{jm}\tilde{\beta}_{jm}\right).$$

The Hamiltonian $\boldsymbol{H}_{\mathrm{TSP}}$ governs dynamics of the system of independent thermal quasiparticles. The ground state of this system is the thermal vacuum state $|0(T)\rangle$ defined as follows:

$$\beta_{jm}|0(T)\rangle = \tilde{\beta}_{jm}|0(T)\rangle = 0.$$

The interaction of thermal quasiparticles is given by the term

$$\boldsymbol{H}_{qph} = -\frac{1}{2}\sum_{\lambda\mu}\sum_{\tau,\rho=\pm1}\left(\chi_0^{(\lambda)} + \chi_1^{(\lambda)}\right)\left[M^{+}_{\lambda\mu}(\tau)M_{\lambda\mu}(\rho\tau) - \widetilde{M}^{+}_{\lambda\mu}(\tau)\widetilde{M}_{\lambda\mu}(\rho\tau)\right]. \qquad (3)$$

In terms of thermal quasiparticles the operator $M^{+}_{\lambda\mu}(\tau)$ reads

$$M^{+}_{\lambda\mu}(\tau) = -\frac{1}{\sqrt{2\lambda+1}}\sum_{j_1 j_2} f^{(\lambda)}_{j_1 j_2}\left[\sqrt{1\text{-}n_{j_1}}\sqrt{n_{j_2}}\,[\beta^{+}_{j_1}\tilde{\beta}^{+}_{\overline{j_2}}]_{\lambda\mu} + \sqrt{n_{j_1}}\sqrt{1-n_{j_2}}\,[\tilde{\beta}_{j_1}\tilde{\beta}_{\overline{j_2}}]_{\lambda\mu} + \right.$$

[1] General formulae can be found in [19, 20, 26].

$$+\sqrt{1-n_{j_1}}\sqrt{1-n_{j_2}}[\beta^+_{j_1}\beta_{\overline{j_2}}]_{\lambda\mu}+\sqrt{n_{j_1}}\sqrt{n_{j_2}}[\widetilde{\beta}_{j_1}\widetilde{\beta}^+_{\overline{j_2}}]_{\lambda\mu}] \quad (4)$$

The square brackets $[\ ]_{\lambda\mu}$ stand for the coupling of single-particle angular momenta j_1 and j_2 to the sum angular momentum λ. The bar over lower index $\overline{j}$ denotes the time reversal state. The value $f^{(\lambda)}_{j_1j_2}$ is a reduced single particle matrix element of the operator $M_{\lambda\mu}$.

The thermal RPA equations are obtained with the phonon operator of the following type:

$$Q^+_{\lambda\mu i}=\sum_{j_1j_2}\eta^{\lambda i}_{j_1j_2}[\beta^+_{j_1}\widetilde{\beta}^+_{\overline{j_2}}]_{\lambda,\mu}+(-1)^{\lambda-\mu}\varsigma^{\lambda i}_{j_1j_2}[\beta_{j_1}\widetilde{\beta}_{\overline{j_2}}]_{\lambda,-\mu}, \quad (5)$$

under the assumption that the ground state of a hot nucleus is the thermal phonon vacuum state $|\Psi_0(T)\rangle$, i.e. $Q_{\lambda\mu}|\Psi_0(T)\rangle=0$. Moreover, the phonon operators (5) are treated as bosonic ones, which mean that the structure of the thermal phonon vacuum state does not deviate strongly from that of the thermal quasiparticle vacuum state $|0(T)\rangle$.

The equation for thermal phonon energies $\omega_{\lambda i}$ is

$$[X_n(\omega)+X_p(\omega)](\chi_0^{(\lambda)}+\chi_1^{(\lambda)})-4\chi_0^{(\lambda)}\chi_1^{(\lambda)}X_n(\omega)X_p(\omega)=1, \quad (6)$$

where

$$X_\tau(\omega)=\frac{1}{2\lambda+1}\sum_{j_1j_2}\frac{\left(f^{(\lambda)}_{j_1j_2}\right)^2(n_{j_1}-n_{j_2})(E_{j_1}-E_{j_2})}{(E_{j_1}-E_{j_2})^2-\omega^2}.$$

The amplitudes $\eta^{\lambda i}_{j_1j_2}$ and $\zeta^{\lambda i}_{j_1j_2}$ are

$$\eta^{\lambda i}_{j_1j_2}=-\sqrt{\frac{1}{2N^{\lambda i}_\tau}}\frac{f^{(\lambda)}_{j_1j_2}\sqrt{1-n_{j_1}}\sqrt{n_{j_2}}}{(E_{j_1}-E_{j_2})-\omega_{\lambda i}},\qquad \zeta^{\lambda i}_{j_1j_2}=-\sqrt{\frac{1}{2N^{\lambda i}_\tau}}\frac{f^{(\lambda)}_{j_1j_2}\sqrt{1-n_{j_1}}\sqrt{n_{j_2}}}{(E_{j_1}-E_{j_2})+\omega_{\lambda i}} \quad (7)$$

In Eq. (7), the value $1/\sqrt{2N^{\lambda i}_\tau}$ is a normalization factor of thermal phonon wave function. Note, that values $N^{\lambda i}_n$ and $N^{\lambda i}_p$ are different.

In contrast with RPA at T = 0 the solutions of (6) with negative energies have physical meanings. They correspond to tilde-phonon states $\widetilde{Q}^+_{\lambda i}|\Psi_0(T)\rangle$.

Coupling of thermal phonon

In terms of TRPA phonons and thermal quasiparticles the thermal Hamiltonian reads

$$\boldsymbol{H}=\sum_{\lambda\mu i}\omega_{\lambda i}\left(Q^+_{\lambda\mu i}Q_{\lambda\mu i}+\tilde{Q}^+_{\lambda\mu i}Q_{\lambda\mu i}\right)- \quad (8)$$

$$-\frac{1}{2\sqrt{2}}\sum_{\lambda\mu i}\sum_\tau\sum_{j_1j_2}\frac{f^{(\lambda)}_{j_1j_2}}{\sqrt{N^{\lambda i}_\tau}}\{\left((-1)^{\lambda-\mu}Q^+_{\lambda\mu i}+Q_{\lambda-\mu i}\right)B(j_1j_2;\ \lambda-\mu)-\left((-1)^{\lambda-\mu}\widetilde{Q}^+_{\lambda\mu i}+\widetilde{Q}_{\lambda-\mu i}\right)\widetilde{B}(j_1j_2;\ \lambda-\mu)+\text{h.c.}\},$$

where

$$B(j_1j_2;\ \lambda\text{-}\mu)=\sqrt{1-n_{j1}}\sqrt{1-n_{j2}}[\beta^+_{j1}\beta_{\overline{j2}}]_{\lambda\mu}+\sqrt{n_{j1}}\sqrt{n_{j2}}[\tilde{\beta}_{j1}\tilde{\beta}^+_{\overline{j2}}]_{\lambda\mu}$$

The item of the thermal Hamiltonian (8) containing the operators $B(j_1j_2;\lambda\mu)$ and its tilde-counterpart is responsible for the interaction of thermal quasiparticles with thermal phonons. It mixes a one-phonon state with more complex configurations.

To describe the fragmentation of thermal phonons we use the variational method with a trial wave function of the form

$$|\Psi_\nu(JM)\rangle = \Big\{ \sum_i R_i(J\nu) Q^+_{JMi} + \sum_{\lambda_1\lambda_2 i_1 i_2} P^{\lambda_1 i_1}_{\lambda_2 i_2}(J\nu) \big[Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2} \big]_{JM} \Big\} |\Psi_0(T)\rangle . \qquad (9)$$

The secular TRPA equation for excited state energies is

$$det \left| (\omega_{Ji} - \eta_{J\nu})\delta_{ii'} - \frac{1}{2} \sum_{\lambda_1 i_1 \lambda_2 i_2} \frac{U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)\, U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji')}{\omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2} - \eta_{J\nu}} \right| = 0 \qquad (10)$$

Functions $U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)$ are coupling matrix elements between one- and two-phonon thermal states. They are proportional to sums of different bilinear combinations of TRPA-phonon amplitudes $\eta^{\lambda i}_{j_1 j_2}$ and $\zeta^{\lambda i}_{j_1 j_2}$. The corresponding formulae can be found in [20, 26].

Within the outlined formalism the calculation procedure is the following: At a given value of T we calculate the energies and structures of TRPA phonons (Eqs. (6) and (7); then all necessary phonon coupling matrix elements $U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)$ are calculated. Then, to calculate the E1 strength function $b(\mathrm{E1},\eta)$ at $\mathrm{T} \neq 0$ taking into account a fragmentation of thermal one-phonon dipole states, we explore the well-known strength function method using the weight function of the Lorentz type [22, 23].

Numerical results

We calculate the E1-strength functions for $0 \leq \mathrm{T} \leq 3$ MeV in ^{120}Sn and ^{208}Pb nuclei. All model parameters (mean field potentials, pairing constants, coupling constants of separable interactions etc) are fixed in accordance with the usual QPM procedure [21, 23], i.e., by the use of experimental data on the energies of low-lying vibrational states and giant resonances at T = 0. As a nuclear mean field the phenomenological Woods-Saxon potential is explored. The single-particle basis consists of all bound states and several quasibound ones with relatively small escape widths.

Only multipole-multipole particle-hole interactions with $1 \leq \lambda \leq 7$ are included in the Hamiltonian. A radial form factor of the separable multipole interaction has the form $\mathrm{R(r)} = \mathrm{dU/dr}$, where U is the central part of the Woods-Saxon potential. The coupling constant of the isoscalar dipole-dipole interaction is adjusted at every value of T to make the energy of the spurious 1^--state zero in the TRPA calculations.

Let us note, that in the thermal RPA (see, [26]) we get qualitatively the same results as many other authors (see, e.g., [27]). When temperature increases, only some minor redistribution of the E1 strength between different one-phonon 1^--states takes place. The energy centroids and Landau widths of GDR in both the nuclei almost do not change with T.

The smearing parameter Δ in the Lorentz weight function is taken to be equal to1 MeV. As a quantitative measure of the GDR spreading width we use a variance σ_{th} of the E1-strength distribution calculated with the following formula:

$$\sigma_{\mathrm{th}} = \sqrt{\frac{m_2}{m_0} - \left(\frac{m_1}{m_0}\right)^2} ,$$

where m_k (k = 0, 1, 2) is the kth energy moment of the E1 strength function defined as follows:

$$m_k = \int_{E_{min}}^{E_{max}} \eta^k\, b(\mathrm{E1},\eta) \mathrm{d}\eta .$$

The calculated T-dependence of σ_{th} together with the values of Γ_{exp} in ^{120}Sn and ^{208}Pb nuclei are displayed in Figs. 1 and 2, respectively.

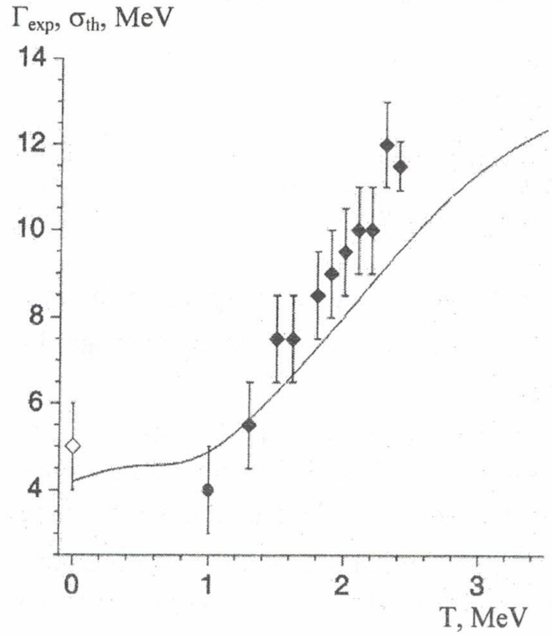

Fig. 1. Temperature dependence of the experimental GDR width Γ_{exp} and the variance σ_{th} of the theoretical E1 strength function in ^{120}Sn. Full diamonds, the revised experimental data from [6]; full circle, the data from [11]; open diamond, the data from [28].

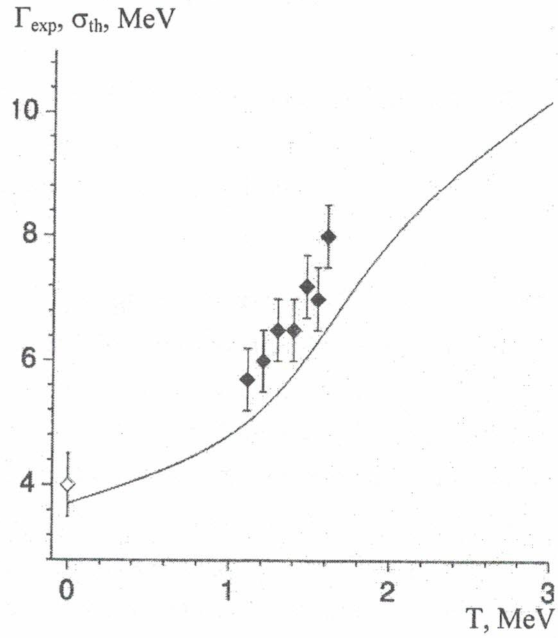

Fig. 2. Temperature dependence of the experimental GDR width Γ_{exp} and the dispersion σ_{th} of the theoretical E1 strength function in ^{208}Pb. Full diamonds, the revised experimental data fro [6], open diamond, the data from [28].

Let us note that direct numerical comparison of σ_{th} with Γ_{exp} cannot be justified because Γ_{exp} values are related to the so-called half-width of the experimental E1 strength function, which is supposed to be of the Lorentz shape [28], but not to a variance. Moreover, the experimental data are related to the total width whereas $\Gamma^{\downarrow}$ is only a part of it. Thus, displaying of σ_{th} and Γ_{exp} on the same figures has only demonstrative sense.

The most distinctive feature of the theoretical curves is that they show increase in the variance and thus in the spreading GDR width with temperature. Moreover, $\sigma_{th} \approx \Gamma_{exp}$ at T = 0 and the temperature dependences of σ_{th} and Γ_{exp} appears to be quite similar.

Discussion

Our results on the behavior of $\Gamma^{\downarrow}(T)$ qualitatively agree with that of [15] and [16, 17] but are in contradiction with those of [12, 14]. To understand why in our approach the value σ_{th} increases with temperature, we analyze the matrix elements of a phonon-phonon coupling $U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)$ and found a strong effect of a coupling of GDR phonons with few very low-lying thermal phonons appearing in the phonon spectrum only at $T \neq 0$ due to the thermal occupation factors. These states correspond to very low-lying poles $E_{j_1} - E_{j_2}$ of the TRPA secular equation (6). One amplitude $\eta^{\lambda i}_{j_1 j_2}$ dominates the phonon wave function, i.e. these phonons are noncollective and of the p-p or h-h type.

In total this conclusion agrees with the results of [15] although in that paper a special role of these low-lying p-p (h-h) phonons was not definitely pointed out. It seems that in [12] the noncollective thermal RPA excitations have been ignored (the same statement can be found in [15]).

Some questions concerning a dependence of our numerical results on parameters still remain. An appearance of low-lying p-p (h-h) states is dependent on the parameters of the mean field. Probably, the use of a phenomenological Saxon-Woods potential gives the upper limit for the role of these low-lying p-p (h-h) states because the density of single-particle states near the Fermi level is the largest one in this potential. We guess that, e.g., with the mean field calculating by the Hartree-Fock method with a density-dependent effective interaction like the Skyrme one, the influence of these phonons on $\Gamma^{\downarrow}$ will be weaker.

One more ingredient directly affecting the calculated value of σ_{th} is the size of the two-phonon state space. To diminish the effect of the Pauli principle violation in these states and to avoid the overcompleteness of the two-phonon basis, we include in the trial wave function (9) only two-phonon configurations combining two collective or one collective – one noncollective phonons. Since the low-lying collective states dissolve with increasing temperature there is no clear cut separation between collective and noncollective states. At finite temperature the fragmentation of GDR appears to be noticeably dependent on the definition of "phonon collectivity" (see, [26]). Nevertheless, a general trend of the thermal behavior of the spreading GDR width is saved.

There is one more interesting difference between our approach and that of [12, 14, 15]. The difference has already been pointed out in [20] and now we would like to discuss it in more detail. In [12, 15], the GDR width depends on thermal occupation numbers of two types – the Fermi - Dirac and the Bose - Einstein occupation numbers. The appearance of Bose occupation factors is a consequence of treating phonons as bosons when the temperature dependent Green's function of a single phonon is introduced.

In the QPM-TFD approach to the phonon-phonon interaction at $T \neq 0$ on can not find the thermal bosonic occupation numbers. The reason is evident: We start with the model Hamiltonian written in terms of nucleonic (i.e. fermionic) variables. The thermal occupation numbers appear in the game when we make the thermal Bogoliubov rotation (2) and thus produce the thermal Fock space. All further manipulations explore these "heated" fermions and there is no room for the appearance of bosonic occupation numbers. Our thermal Hamiltonian in its final form (8) is the Hamiltonian of interacting phonons built from "heated" quasiparticles but the phonon system itself is not heated in the sense that there is no thermal smearing of phonons over their energy levels.

This corresponds to a transparent phenomenological picture: when one heats a nucleus putting there a good piece of energy, a nucleonic motion is changed and due to this the properties of a nuclear surface are changed. However, this doesn’t mean that nuclear surface vibrations are heated themselves.

In [15], the authors start just with the Hamiltonian of the interacting TRPA phonons implying, as an obvious fact, that the phonon system has the same temperature T as the underlying fermions forming the thermal phonons. In our opinion this is an additional assumption which has to be justified. Similarly, in [14] from the beginning a nucleus is treated as a system of phonons and quasiparticles. However, since phonons and quasiparticles are considered as some "initial" ingredients, the structure of phonons has to be as it is in a cold nucleus and cannot be changed by heating the system. Thus, they cannot satisfy the thermal RPA equation.

The point is that quasiparticles and phonons are not independent variables in a nucleus. The phonon is a coherent superposition of two-quasiparticle components. So, starting with the model Hamiltonian given in terms of nucleonic degrees of freedom one has to make a mapping of pure fermionic states to a subspace consisting of ideal "quasiparticle" and "bosonic" elementary modes.

In this regard, Hatsuda [25] discussed already two ways to consider a hot nucleus in the framework of TFD. The first is to make a mapping of the initial Hamiltonian and the initial pure fermionic Fock space of a cold system (nucleus) and only after this to thermalize a system in question. For the approach presented here it means that degrees of freedom should be doubled for the quasiparticle-phonon image of the Hamiltonian (1) at T = 0. Then one gets the thermal Hamiltonian with both the types of thermal occupation numbers and, consequently, the GDR width also should depend on them. However, Hatsuda [25] has also shown taking the Lipkin model as an example that "thermalizing" of the bosonic image of the initial fermionic Hamiltonian one cannot derive in the leading order the TRPA equations for these bosons.

The second way is just our present way: while heating we treat a nucleus as a system of fermions and only after this we project or transform the original nucleonic degrees of freedom to more convenient ones (bosonic or bosonic + fermionic).

We would like to stress that the problem how to treat a thermalized nucleus in terms of quasiparticles and phonons is not so trivial as it may seem at the first glance. It is in intimate correspondence with a proper choice of physically important degrees of freedom and their consistent mapping which has to comply with the particle statistic requirements [29].

Conclusions

A temperature dependence of the fragmentation of a giant dipole resonance has been studied within the quasiparticle – phonon model extended to finite temperatures within the thermo field dynamics. The increase of the variance of the theoretical E1 strength function with temperature in the range $0 < T \leq 3$ MeV has been found.

Moreover, we draw attention to the problem of a proper selection of relevant nuclear degrees of freedom to describe a mode-mode coupling in a hot nucleus. To our knowledge, this aspect of nuclear theory was overlooked before.

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ЗАЛЕЖНІСТЬ ШИРИНИ ГІГАНТСЬКОГО ДИПОЛЬНОГО РЕЗОНАНСУ ВІД ТЕМПЕРАТУРИ

Залежність спредової ширини $\Gamma^{\downarrow}$ гігантського дипольного резонансу від температури досліджено в рамках квазічастково-фононної моделі ядра, узагальненої на ненульові температури за допомогою формалізму термопольової динаміки. Чисельні розрахунки, виконані для компаунд-ядер ^{120}Sn та ^{208}Pb, указують на швидкий ріст з температурної ширини $\Gamma^{\downarrow}$, що якісно узгоджується з експериментом. Обговорюються причини цього явища, а також зв'язок запропонованого підходу з іншими теоріями.

ЗАВИСИМОСТЬ ШИРИНЫ ГИГАНТСКОГО ДИПОЛЬНОГО РЕЗОНАНСА ОТ ТЕМПЕРАТУРЫ

А. И. Вдовин, А. Н. Стороженко

Зависимость спредовой ширины $\Gamma^{\downarrow}$ гигантского дипольного резонанса от температуры исследована в рамках квазичастично-фононной модели ядра, обобщенной на ненулевые температуры с помощью формализма термополевой динамики. Численные расчеты, выполненные для компаунд-ядер ^{120}Sn и ^{208}Pb, указывают на быстрый рост с температурой ширины $\Gamma^{\downarrow}$, что качественно согласуется с экспериментом. Обсуждаются причины этого явления, а также связь предложенного подхода с другими теориями.

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