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EQUATIONS OF THE FINITE TEMPERATURE QUASIPARTICLE-PHONON NUCLEAR MODEL FOR HOT SPHERICAL ODD NUCLEI

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1. Introduction

Recently very excited nuclei can be produced in heavy-ion collisions and their properties can be studied experimentally as functions of temperature. Experimental evidences have suggested that a nuclear system is thermally equilibrated rapidly after formation of a compound nucleus from deep-inelastic nuclear collisions andheavy-ion fusion [1]. The large nuclear state density equally populates individual highly-excited states. Consequently, average properties of the system are usually measured. It is why the statistical extension of different microscopic approaches can be appropriate to be applied to study collective states at finite temperature. The first step in this extension is the finite temperature RPA (FT-RPA)[2-7]. To describe the fragmentation of collective states one should include the coupling to more complicated configurations. Thus, the investigations of the damping of giant resonances (GR) in hot nuclei have led to the extension of the Nuclear Field Theory (NFT)[8] to finite temperature based on the finite temperature Matsubara formalism [9], the second FT-RPA[10] and the finite temperature Quasiparticle-Phonon Nuclear Model (FT-QPNM)[11,12]. The latter is an extension of the QPNM[13] for cold even-even spherical nuclei to finite temperature where the damping of GR is understood as the fragmentation of thermal one-phonon states due to the coupling to (2p-2h) configurations. The structures of thermal one-phonons are calculated in the FT-RPA and the FT-QPNM Hamiltonian for even-even nuclei can be expressed in terms of these FT-RPA operators. Thus, we have obtained in Refs.[11,12] the set of basic FT-QPNM equations for hot even-even spherical nuclei. The calculations performed in our approach as well as in the NFT have shown that the damping of the giant dipole resonances (GDR) in hot even-even spherical nuclei is weakly dependent on temperature.

In odd A nuclei the fragmentation of single-quasiparticle states is mainly governed in the QPNM by the coupling to the "quasiparticle ⊗ phonon" states[13]. The properties of excited states in odd A nuclei at zero temperature have been studied in detail in many works within the scope of the QPNM during the last decade /See [13] and refs. therein/. The extension of the QPNM to finite temperature for hot spherical odd-A nuclei is however absent so far.

The aim of the present effort is to make the first step towards the generalization of the QPNM to finite temperature for describing hot spherical odd-A nuclei. We shall derive the set of the FT-QPNM equations in hot spherical odd nuclei for the simplest case with configurations not more complicated than "quasiparticle ⊗ phonon" ones.

2. Model Hamiltonian. Thermal vacuum and statistical ensemble

Consider the QPNM (FT-QPNM) Hamiltonian in the form[13]

$$H = H_\alpha + H_Q + H_{\alpha Q} . \qquad (1)$$

where H_α describes the motion of independent quasiparticles

$$H_\alpha = \sum_{jm} \varepsilon_j \alpha^+_{jm} \alpha_{jm} . \qquad (2)$$

H_Q stands for the independent phonon part of the Hamiltonian

$$H_Q = -\frac{1}{4} \sum_{\lambda\mu i i' \tau} \frac{X^{\lambda i}(\tau) + X^{\lambda i'}(\tau)}{\left[\mathcal{Y}^{\lambda i}_\tau \mathcal{Y}^{\lambda i'}_\tau \right]^{1/2}} Q^+_{\lambda\mu i} Q_{\lambda\mu i'} \qquad (3)$$

and $H_{\alpha Q}$ corresponds to the couplings of quasiparticles and phonons

$$H_{\alpha Q} = -\sum_{\lambda\mu i} \sum_{jj'}^{\tau} \Gamma^{\lambda i}_{jj'}(\tau) \{ [Q^+_{\lambda\mu i} + Q_{\lambda\tilde{\mu} i}] B_{\lambda\tilde{\mu}}(jj') + h.c. \} . \qquad (4)$$

In Eqs.(2)-(4) ε_j are the quasiparticle energies; α^+_{jm} and α_{jm} are the creation and annihilation quasiparticle operators, respectively. The phonon operators $Q^+_{\lambda\mu i}$ and $Q_{\lambda\mu i}$ are defined as [13]

$$Q^+_{\lambda\mu i} = \frac{1}{2} \sum_{jj'} \{ \psi^{\lambda i}_{jj'} A^+_{\lambda\mu}(jj') - \phi^{\lambda i}_{jj'} A_{\lambda\tilde{\mu}}(jj') \}$$

$$Q_{\lambda\mu i} = (Q^+_{\lambda\mu i})^+ . \qquad (5)$$

The pair creation (annihilation) A^+, A and scattering B, B^+ quasiparticle operators are

$$A^+_{\lambda\mu}(jj') = \sum_{mm'} \langle jmj'm' | \lambda\mu \rangle \alpha^+_{jm} \alpha^+_{j'm'} \equiv [\alpha^+_j \otimes \alpha^+_{j'}]_{\lambda\mu} \qquad (6)$$

$$A_{\lambda\mu}(jj') = \sum_{mm'} \langle jmj'm' | \lambda\mu \rangle \alpha_{j'm'} \alpha_{jm} \equiv [\alpha_{j'} \otimes \alpha_j]_{\lambda\mu} \qquad (7)$$

$$B_{\lambda\mu}(jj') = -\sum_{mm'} \langle jmj'm' | \lambda\mu \rangle \alpha^+_{jm} \alpha_{j'\tilde{m}'} \equiv -[\alpha^+_j \otimes \tilde{\alpha}_{j'}]_{\lambda\mu} \qquad (8)$$

$$B^+_{\lambda\mu}(jj') = (-)^{j-j'-\lambda} B_{\lambda\tilde{\mu}}(j'j), \qquad (9)$$

where the time inversed notation $\mathcal{O}_{\lambda\tilde{\mu}} \equiv (-)^{\lambda-\mu} \mathcal{O}_{\lambda-\mu}$ for an arbitrary operator $\mathcal{O}_{\lambda\mu}$ is used.

The notation $\tau = \{ n, p \}$ denotes neutron and proton components. The change $\tau \leftrightarrow -\tau$ corresponds to $n \leftrightarrow p$. However, in further consideration we will omit τ in the formulae for simplicity.

At zero temperature the functions $X^{\lambda i}$ and $\mathcal{Y}^{\lambda i}$ in Eq.(3) are calculated from the RPA equations. Their explicit expressions have been given in Refs. /13/. At finite temperature $X^{\lambda i}$ and $\mathcal{Y}^{\lambda i}$ have the form defined from the FT-RPA equations, which can be found in Refs. /6,12/.

The abbreviation $\Gamma^{\lambda i}_{jj'}$ denotes /13/

$$\Gamma^{\lambda i}_{jj'} = \left(8\mathcal{Y}^{\lambda i}\right)^{-1/2} f^{(\lambda)}_{jj'} v^{(-)}_{jj'} , \tag{10}$$

where $f^{(\lambda)}_{jj'}$ are the reduced matrix elements for single particle operators generating excitations.
The coefficient $v^{(-)}_{jj'}$ in Eq.(10) is a combination of the Bogolubov coefficients u_j , v_j . It reads $v^{(-)}_{jj'} = u_j u_{j'} - v_j v_{j'}$.
In fact, in the Hamiltonian (1) we have neglected the term containing the combinations $\sim BB$ of operators B from Eqs.(8)-(9). At zero temperature the estimation performed in Ref. /14/ has shown the negligibly small contribution of this term to the one-phonon energies calculated in the RPA. At finite temperature a part of this term leading to new poles is included to the definition of the FT-RPA phonons and therefore renormalizes the thermal one-phonon energies.

At zero temperature in general the ground state of odd--nuclei is taken as the mixed quasiparticle $\otimes$ phonon vacuum $|0\rangle$ That means

$$|0\rangle = |0\rangle_\alpha \otimes |0\rangle_Q , \tag{11}$$

where $|0\rangle_\alpha$ is the quasiparticle vacuum

$$\alpha_{jm} |0\rangle_\alpha = 0 \tag{12}$$

and $|0\rangle_Q$ is the phonon vacuum

$$Q_{\lambda\mu i} |0\rangle_Q = 0 . \tag{13}$$

At finite temperature a statistical ensembe of quasiparticle and phonon excitations is obtained. The vacuum (11) can no longer serve as a reference state to define the normal product. Instead of it the thermal vacuum $|0,\beta\rangle$ with $\beta = T^{-1}$ the inversed temperature must be used. Its explicit form is given in Ref /15/.

It has been shown in /15/ that the statistical ensemble average of an operator $\hat{\mathcal{O}}$ can be expressed as the vacuum expectation value

$$\langle\hat{\mathcal{O}}\rangle = \langle 0,\beta|\hat{\mathcal{O}}\otimes\hat{1}|0,\beta\rangle , \tag{14}$$

where

$$\langle\hat{\mathcal{O}}\rangle \equiv Tr(W\hat{\mathcal{O}})$$
$$W \equiv exp(-\beta H^{eff})/Tr\left[exp(-\beta H^{eff})\right] \tag{15}$$

is the ensemble average of operator $\hat{\mathcal{O}}$. Hereafter the ensemble average (15) is always understood at finite temperature.

3. Equations for finite temperature Green functions in hot spherical odd nuclei.

As has been mentioned in Sec. 2, the phonon operators (5) have quasiparticle (fermion) structure. The exact commutation relations for operators (5) as well as between operators (5) and quasiparticle operators have been given in many works devoted to the QPNM. Taking into account these commutators is equivalent to the exact inclusion of the Pauli principle (See, e.g. /16-18/). The effect of the Pauli principle has been the subject of many papers within the framework of the QPNM. It belongs to the higher order corrections in the perturbative expansion of the theory. Therefore, in the present work, for simplicity, we will neglect it supposing the quasi-boson structure for phonons everywhere if any special note is not made. Thus, we have approximately /13/

$$[Q_{\lambda\mu i}, Q^+_{\lambda'\mu' i'}] \approx \delta_{\lambda\lambda'}\delta_{\mu\mu'}\delta_{ii'} \tag{16}$$

$$[Q_{\lambda\mu i}, Q_{\lambda'\mu' i'}] \approx [Q^{+}_{\lambda\mu i}, Q^{+}_{\lambda'\mu' i'}] \approx 0, \qquad (17)$$

which is the well-known quasiboson approximation. With these "ideal bosons" $Q_{\lambda\mu i}$, $Q^{+}_{\lambda'\mu' i'}$ we also have

$$[\alpha_{jm}, Q_{\lambda\mu i}] \approx 0 \quad ; \quad [\alpha^{+}_{jm}, Q_{\lambda\mu i}] \approx 0 \qquad (18)$$

and consequently

$$[B_{\lambda\mu}(jj'), Q_{\lambda'\mu i}] \approx 0 \quad ; \quad [B_{\lambda\mu}(jj'), Q^{+}_{\lambda'\mu i}] \approx 0 . \qquad (19)$$

The quasiparticle operators α^{+}_{jm} and α_{jm} satisfy the usual anticommutation relations for fermions.

The commutation relation between the quasiparticle operator and the scattering quasiparticle operator B from Eqs. (8), (9) is

$$[\alpha_{jm}, B_{\lambda\mu}(j_1 j_2)] = -\delta_{jj_1} \sum_{m_2} \langle jm j_2 m_2 | \lambda\mu \rangle \alpha_{j_2 \tilde{m}_2} . \qquad (20)$$

We now use Eqs.(16)-(20) for deriving the set of equations for the finite temperature Green functions of interest.

We define the two-temporal Green functions at finite temperature that represent:

a) the processes with one propagating unpaired (odd) quasiparticle

$$G^{-,+}_{JM;J'M'}(t-t') = -i\theta(t-t')\langle[\alpha_{JM}(t), \alpha^{+}_{J'M'}(t')]\rangle$$
$$\equiv \langle\langle \alpha_{JM}(t); \alpha^{+}_{J'M'}(t') \rangle\rangle . \qquad (21)$$

b) the transition between the unpaired quasiparticle and the coupling of a phonon with a quasiparticle

$$G^{--,+}_{JM,\Lambda\Sigma I;J'M'}(t-t') = -i\theta(t-t')\langle[\alpha_{JM}(t)\, Q_{\Lambda\Sigma I}(t), \alpha^{+}_{J'M'}(t')]\rangle$$
$$\equiv \langle\langle \alpha_{JM}(t)\, Q_{\Lambda\Sigma I}(t); \alpha^{+}_{J'M'}(t') \rangle\rangle , \qquad (22)$$

c) the quasiparticle "scattering" with "emitting" an intermediate phonon

$$G^{-+,+}_{JM,\Lambda\Sigma I;J'M'}(t-t') = -i\theta(t-t')\langle[\alpha_{JM}(t)\, Q^{+}_{\Lambda\tilde{\Sigma} I}(t), \alpha^{+}_{J'M'}(t')]\rangle$$
$$\equiv \langle\langle \alpha_{JM}(t)\, Q^{+}_{\Lambda\tilde{\Sigma} I}(t); \alpha^{+}_{J'M'}(t') \rangle\rangle . \qquad (23)$$

Hereafter, the fixed indices are denoted by capital Latin letters $JM,\ldots$ for quasiparticles, and by capital Greek letters $\Lambda\Sigma I,\ldots$ for phonons, while the corresponding small letters ($jm,\ldots$, $\lambda\mu i,\ldots$) are used in the sums.

The diagrammatic representation the processes (21)-(23) is shown in Fig.1, that illustrates the Eq.(33) (see below).

Functions (21)-(23) include all the forward processes with the contribution of quasiparticle and quasiparticle phonon configurations corresponding to the Hamiltonian (1)-(4).

By applying the standard thermodynamic Green function method /19-20/ we note that due to the RPA (FT-RPA) solutions the parts (2) and (3) of the Hamiltonian (1) can be unified in the harmonic part

$$H_h \equiv H_\alpha + H_Q = \sum_{\lambda\mu i} \omega_{\lambda i}\, Q^{+}_{\lambda\mu i} Q_{\lambda\mu i} . \qquad (24)$$

This form is more convenient for calculating the commutators between the Hamiltonian (1) and the phonon operators (5) under the approximation (16)-(20).

a) b)

Fig. 1 Graphs summarized in Eq.(33)

a) The sum $\sum_{\lambda i j} \Gamma^2(Jj\lambda i) \dfrac{1-n_j+\nu_{\lambda i}}{\varepsilon_j+\omega_{\lambda i}-\eta}$

b) The sum $\sum_{\lambda i j} \Gamma^2(Jj\lambda i) \dfrac{n_j+\nu_{\lambda i}}{\varepsilon_j-\omega_{\lambda i}-\eta}$.

By closing the hierarchy with the Green functions (21)-(23), we use the well-known decoupling procedure [19,20]

$$\langle\langle \alpha^+_{jm}(t)\,\alpha_{JM}(t)\,\alpha_{j'm'}(t);\,\alpha^+_{J'M'}(t')\rangle\rangle \approx$$

$$\approx \delta_{jJ}\,\delta_{mM}\,\langle\alpha^+_{JM}\,\alpha_{JM}\rangle\langle\langle\alpha_{j'm'}(t);\,\alpha^+_{J'M'}(t')\rangle\rangle = \delta_{jJ}\,\delta_{mM}\,n_J\,G^{-,+}_{j'm';J'M'}(t-t')$$

$$\langle\langle Q^+_{\lambda\mu i}(t)\,Q_{\Lambda\Sigma I}(t)\,\alpha_{jm}(t);\,\alpha^+_{J'M'}(t')\rangle\rangle \approx \qquad (25)$$

$$\approx \delta_{\lambda\Lambda}\,\delta_{\mu\Sigma}\,\delta_{iI}\,\langle Q^+_{\Lambda\Sigma I}\,Q_{\Lambda\Sigma I}\rangle\langle\langle\alpha_{jm}(t);\,\alpha^+_{J'M'}(t')\rangle\rangle = \delta_{\lambda\Lambda}\,\delta_{\mu\Sigma}\,\delta_{iI}\,\nu_{\Lambda I}\,G^{-,+}_{jm;J'M'}(t-t'),$$

where

$$n_J = [\,exp(\beta\varepsilon_J)+1\,]^{-1} \qquad (26)$$

$$\nu_{\Lambda I} = [\,exp(\beta\omega_{\Lambda I})-1\,]^{-1} \qquad (27)$$

are respectively the quasiparticle and phonon occupation numbers.

After a rather lengthy but straightforward algebra, we obtain the set of equations for the two temporal Green functions (21)-(23). Taking their Fourier images we obtain for them the set of equations

$$\begin{cases} (\varepsilon_J-\eta)\,G^{-,+}_{JM;J'M'}(\eta) - 2\sum\limits_{\lambda\mu i}\sum\limits_{jm}\Gamma^{\lambda i}_{Jj}\langle JMj\tilde{m}|\lambda\mu\rangle[\,G^{--,+}_{jm,\lambda\mu i;J'M'}(\eta)+G^{-+,+}_{jm,\lambda\mu i;J'M'}(\eta)\,] = \\ \qquad\qquad = \frac{1}{2\pi}\,\delta_{JJ'}\,\delta_{MM'} \qquad (28) \\ (\varepsilon_J+\omega_{\Lambda I}-\eta)\,G^{--,+}_{JM,\Lambda\Sigma I;J'M'}(\eta) - 2\sum\limits_{jm}\Gamma^{\Lambda I}_{Jj}\langle JMj\tilde{m}|\Lambda\tilde{\Sigma}\rangle(1-n_J+\nu_{\Lambda I})\,G^{-,+}_{jm;J'M'}(\eta) = 0 \qquad (29) \\ (\varepsilon_J-\omega_{\Lambda I}-\eta)\,G^{-+,+}_{JM,\Lambda\Sigma I;J'M'}(\eta) - 2\sum\limits_{jm}\Gamma^{\Lambda I}_{Jj}\langle JMj\tilde{m}|\Lambda\tilde{\Sigma}\rangle(n_J+\nu_{\Lambda I})\,G^{-,+}_{jm;J'M'}(\eta) = 0. \qquad (30) \end{cases}$$

Expressing $G^{--,+}$ and $G^{-+,+}$ from Eqs.(29) and (30) of this set through $G^{-,+}$ and inserting them into Eq.(28), we obtain one equation with one unknown function

$$\left[(\varepsilon_J-\eta)-\sum_{\lambda ij}\Gamma^2(Jj\lambda i)\left(\frac{1-n_j+\nu_{\lambda i}}{\varepsilon_j+\omega_{\lambda i}-\eta}+\frac{n_j+\nu_{\lambda i}}{\varepsilon_j-\omega_{\lambda i}-\eta}\right)\right]G^{-,-}_{JM;J'M'}(\eta)=\frac{1}{2\pi}\,\delta_{JJ'}\,\delta_{MM'}, \qquad (31)$$

where we use the notation [13]

$$\Gamma(Jj\lambda i) \equiv 2\sqrt{\frac{2\lambda+1}{2J+1}}\;\Gamma^{\lambda i}_{Jj}\,. \qquad (32)$$

We note that the Clebsh-Gordon coefficients have been suppressed due to their symmetry properties to give Γ^2 in the sum in Eq.(31). In the homogeneous case (J, M and J', M' are arbitrary) we obtain from Eq.(31) the secular equation for finding the energy η in hot spherical odd nuclei. It reads

$$\varepsilon_J-\eta-\sum_{\lambda ij}\Gamma^2(Jj\lambda i)\left(\frac{1-n_j+\nu_{\lambda i}}{\varepsilon_j+\omega_{\lambda i}-\eta}+\frac{n_j+\nu_{\lambda i}}{\varepsilon_j-\omega_{\lambda i}-\eta}\right)=0\,. \qquad (33)$$

It can easily be seen from Eq.(33) that in the case as $T\to 0$ ($n_j\to 0$, $\nu_{\lambda i}\to 0$) Eq.(33) transforms completely into the well-known QPNM secular equation for spherical odd nuclei with the effect of Pauli principle omitted [13,21]. The second term in the sum in Eq.(33) containing the denominator $\sim(\varepsilon_j-\omega_{\lambda i}-\eta)^{-1}$ is stimulated by the processes (23) which take place only at finite temperature. The appearance of this term produces the new poles ($\varepsilon_j-\omega_{\lambda i}$) at finite temperature which can be located near zero. The question is what contribution these new poles gove to the total strength distribution. As we will see in Sect. 7, the results in a schematic model can however shed light on this problem.

We notice that if we have in mind the fermion structure of phonon onerators, we must use instead of (16) and (19) the exact commutation relations. In this case, we can include the

Pauli principle between "quasiparticle $\otimes$ phonon" components, as has done in Ref. /17/ within the QPNM framework. Thus, in the so-called diagonal in $\mathcal{L}$ approximation of taking account the Pauli-proneiple between "quasiparticle $\otimes$ phonon" configurations we would have in Eqs. (31) and (33) $\Gamma^2(Jj\lambda i)[1+\mathcal{L}(Jj\lambda i)]$ instead of $\Gamma^2(Jj\lambda i)$. The denominators are replaced by $[\varepsilon_j+\omega_{\lambda i}+R(Jj\lambda i)-\eta]$ and $[\varepsilon_j-\omega_{\lambda i}-R(Jj\lambda i)-\eta]$, respestively. The factor $\mathcal{L}(Jj\lambda i)=-1$ when the Pauli principle is violated maximally and such "quasiparticle $\otimes$ phonon" components are excluded automatically from the sum over $(Jj\lambda i)$. The shift $R(Jj\lambda i)$ of the poles appears due to the Pauli principle. The explicit expressions for the factor $\mathcal{L}$ and the shift R are given Refs. /13,17/ and we do not repeat them here.

4. One-to-one correspondence to the conception of excitation operators. Diagrammatic representation

The QPNM uses the conception of multicomponent wave functions obtained by acting the defined operator of excitation on the corresponding ground state wave function. In even-even nuclei, this ground state has been taken as the phonon vacuum $|0\rangle_Q$ (13). In odd nuclei, it is the phonon vacuum of the even-even core which serves the wave function for the ground state /13/. As has been shown in Ref. /12/, in the FT-QPNM the thermal phonon vacuum $|0,\beta\rangle_Q$ can serve as the thermal ground state in hot spherical even-even nuclei. The one-to-one correspondence between the Green functions obtained in that case and the phonon components of the FT-QPNM excitation operator in hot spherical even-even nuclei has been established in Ref. /11/. In the present case of hot odd spherical nuclei, an analogous correspondence can also be pointed out.

Thus, if we define the wave functions for excited states above the thermal vacuum $|0,\beta\rangle_Q$ as

$$\Psi_\nu(JM)_T = Q^+_{JM\nu}|0,\beta\rangle_Q , \tag{34}$$

where

$$Q^+_{JM\nu} = C_{J\nu}\{\alpha^+_{JM}+\sum_{\lambda ij}[D_j^{\lambda i}(J\nu)[\alpha^+_j\otimes Q^+_{\lambda i}]_{JM}+F_j^{\lambda i}(J\nu)[\alpha^+_j\otimes\tilde{Q}_{\lambda i}]_{JM}]\}$$

$$Q_{JM\nu}|0,\beta\rangle_Q = {}_Q\langle 0,\beta|Q^+_{JM\nu} = 0 \tag{35}$$

is the excitation operator for hot spherical odd nuclei, then the one-to-one correspondence between the Green functions (21)-(23) and the coefficients C, D, F in Eq.(35) is

$$\begin{aligned} G^{-,+}_{J\nu;J\nu} &\longleftrightarrow C_{J\nu} \\ G^{--,+}_{j\lambda i;J\nu} &\longleftrightarrow C_{J\nu}D_j^{\lambda i}(J\nu) \\ G^{-+,+}_{j\lambda i;J\nu} &\longleftrightarrow C_{J\nu}F_j^{\lambda i}(J\nu). \end{aligned} \tag{36}$$

Due to Eqs. (28)-(31) the Green functions (21)-(23) do not depend upon the z-projection M.

At zero temperature the Green function $G^{-+,+}$ (23), and consequently the coefficient F, vanish. One then obtains from Eqs. (34) and (35) the usual definition for the wave function of excited states with components not more complicated than $[\alpha^+\otimes Q^+]$ in spherical odd nuclei of the QPNM (T=0) /13/.

The orthonormalized condition for the wave functions (34)

$$\Psi^*_\nu(JM)_T\Psi_\nu(JM)_T = 1 \tag{37}$$

leads to the following equation for the coefficients $C_{J\nu}$, $D_j^{\lambda i}(J\nu)$ and $F_j^{\lambda i}(J\nu)$

$$C^2_{J\nu}\{1-n_J+\sum_{\lambda ij}[D_j^{\lambda i}(J\nu)]^2(1-n_j)(1+\nu_{\lambda i})+\sum_{\lambda ij}[F_j^{\lambda i}(J\nu)]^2(1-n_j)\nu_{\lambda i}\}=1 . \tag{38}$$

It is clear that as $T\to 0$ one obtains from Eq.(38) the well known orthonormalized condition for the coefficients $C_{J\nu}$ and $D_j^{\lambda i}(J\nu)$ in cold odd spherical nuclei.

The wave functions for excited states at $T=0$ can be obtained from Eq.(34) by putting $F_j^{\lambda i}(J\nu)=0$. In the case of including the Pauli principle between "quasiparticle ⊗ phonon" components in the diagonal in $\mathcal{L}$ approximation, we have instead of Eq.(38) the condition

$$C^2_{J\nu}\Big\{1-n_J+\sum_{\lambda ij}[1+\mathcal{L}(Jj\lambda i)]\big[(D_j^{\lambda i}(J\nu))^2(1-n_j)(1+\nu_{\lambda i})+ (F_j^{\lambda i}(J\nu))^2(1-n_j)\nu_{\lambda i}\big]\Big\}=1\,, \qquad (39)$$

which transforms in the limit $T\to 0$ into /13,17/

$$C^2_{J\nu}\Big\{1+\sum_{\lambda ij}[D_j^{\lambda i}(J\nu)]^2[1+\mathcal{L}(Jj\lambda i)]\Big\}=1\,. \qquad (40)$$

Let us denote a quasiparticle by a direct line and a phonon by a wavy line. The arrows on these lines indicate the processes of creation or annihilation of quasiparticle and phonon (quasi-hole and phonon-hole). The time always flows from left to right. The point stands for the vertex $\Gamma(Jj\lambda i)$ (32).

These diagrammatic elements allows us to depict the graphs summarized in Eq.(33) as in Fig. 1.
In Fig.1 the part b) appears only at finite temperature. The graphs including the Pauli principle between "quasiparticle⊗phonon" components in the diagonal in $\mathcal{L}$ approximation are presented in Fig.2.
It is clear from these figures that there is a difference from hot even-even spherical nuclei, where the phonon scattering effect appearing at finite temperature is in a higher order in the perturbation expansion as compared to the main processes with creating (annihilating) two intermediate phonons. In fact, in hot odd spherical nuclei the graphs from Figs. 1b and 2b are in the same order with the ones from Figs. 1a and 2a. Therefore one can expect a more noticeable contribution of the

quasiparticle-phonon scattering at finite temperature in hot odd spherical nuclei.

a) b)

Fig. 2 Graphs including the Pauli principle between "quasiparticle ⊗ phonon" components

a) The sum $\sum_{\lambda ij}\Gamma^2(Jj\lambda i)\dfrac{[1+\mathcal{L}(Jj\lambda i)](1-n_j+\nu_{\lambda i})}{\varepsilon_j+\omega_{\lambda i}+R(Jj\lambda i)-\eta}$

b) The sum $\sum_{\lambda ij}\Gamma^2(Jj\lambda i)\dfrac{[1+\mathcal{L}(Jj\lambda i)]\,(n_j+\nu_{\lambda i})}{\varepsilon_j-\omega_{\lambda i}-R(Jj\lambda i)-\eta}$.

5. Formulae for the damping, spectral intensity and the strength functions in hot odd spherical nuclei.

Formulae for the damping and the spectral intensity can be introduced by analogy with the system describing the interaction of electrons with the lattice in metals /19,20,22/.
Thus, from Eq.(33) we obtain

$$\gamma_J(\eta)=\frac{\Delta}{2}\sum_{\lambda ij}\Gamma^2(Jj\lambda i)\Big[\frac{1-n_j+\nu_{\lambda i}}{(\varepsilon_j+\omega_{\lambda i}-\eta)^2+\Delta^2/4}+\frac{n_j+\nu_{\lambda i}}{(\varepsilon_j-\omega_{\lambda i}-\eta)^2+\Delta^2/4}\Big] \qquad (41)$$

for the damping of the elementary excitations in the system of interacting quasiparticles and phonons, and

$$J_J(\eta)=\frac{1}{\pi}\,\frac{\gamma_J(\eta)(\exp(\beta\eta)+1)^{-1}}{[\varepsilon_J-\eta-\mathcal{M}_J(\eta)]^2+\gamma_J^2(\eta)} \qquad (42)$$

with the mass operator

$$\mathcal{M}_J(E)=\sum_{\lambda ij}\Gamma^2(Jj\lambda i)\left(\frac{1-n_j+\nu_{\lambda i}}{\varepsilon_j+\omega_{\lambda i}-E}+\frac{n_j+\nu_{\lambda i}}{\varepsilon_j-\omega_{\lambda i}-E}\right) \tag{43}$$

for the spectral intensity.

Using Eqs.(41)-(43) we derive the strength functions describing the fragmentation of single-quasiparticle states at finite temperature as (See, e.g. /13/).

$$C_J^2(\eta) = \frac{1}{\pi}\,\frac{\frac{\Delta}{2}\left[1+\mathcal{T}_1(\eta)\right]}{\left[\varepsilon_J-\eta-\gamma_1(\eta)\right]^2+\frac{\Delta^2}{4}\left[1+\mathcal{T}_1(\eta)\right]^2}, \tag{44}$$

where in difference with the zero temperature case in /13/ we have

$$\gamma_1(\eta)=\sum_{\lambda ij}\Gamma^2(Jj\lambda i)\left[\frac{(\varepsilon_j+\omega_{\lambda i}-\eta)(1-n_j+\nu_{\lambda i})}{(\varepsilon_j+\omega_{\lambda i}-\eta)^2+\Delta^2/4}+\frac{(\varepsilon_j-\omega_{\lambda i}-\eta)(n_j+\nu_{\lambda i})}{(\varepsilon_j-\omega_{\lambda i}-\eta)^2+\Delta^2/4}\right]$$

$$\mathcal{T}_1(\eta)=\sum_{\lambda ij}\Gamma^2(Jj\lambda i)\left[\frac{1-n_j+\nu_{\lambda i}}{(\varepsilon_j+\omega_{\lambda i}-\eta)^2+\Delta^2/4}+\frac{n_j+\nu_{\lambda i}}{(\varepsilon_j-\omega_{\lambda i}-\eta)^2+\Delta^2/4}\right].$$

The strength functions for the fragmentation of "quasiparticle ⊗ phonon" states are defined as /13/

$$D^2(\eta)=\frac{\Delta}{2\pi}\sum_\nu\left[C_{J\nu}D_j^{\lambda i}(J\nu)\right]^2\left[(\eta-\eta_{J\nu})^2+\Delta^2/4\right]^{-1} \tag{45}$$

for the process (22) depicted in Fig. 1a, and

$$F^2(\eta) = \frac{\Delta}{2\pi}\sum_\nu\left[C_{J\nu}F_j^{\lambda i}(J\nu)\right]^2\left[(\eta-\eta_{J\nu})^2+\Delta^2/4\right]^{-1} \tag{46}$$

for the process (23) depicted in Fig. 1b.

The functions in the r.h.s. of Eqs. (45) and (46) have the form

$$\left[D_j^{\lambda i}(J\nu)\right]^2=\Gamma^2(Jj\lambda i)(1-n_j+\nu_{\lambda i})^2/(\varepsilon_j+\omega_{\lambda i}-\eta_{J\nu})^2 \tag{47}$$

$$\left[F_j^{\lambda i}(J\nu)\right]^2=\Gamma^2(Jj\lambda i)(n_j+\nu_{\lambda i})^2/(\varepsilon_j-\omega_{\lambda i}-\eta_{J\nu})^2. \tag{48}$$

The function $C^2_{J\nu}$ at the solution points of Eq.(33) is defined as /13,23/

$$C^2_{J\nu} = -\left[d\mathcal{F}(\eta)/d\eta\Big|_{\eta=\eta_{J\nu}}\right]^{-1}, \tag{49}$$

where $\mathcal{F}(\eta)$ is the l.h.s. of Eq.(33).

We can also define the total strength function describing the fragmentation in the "quasiparticle ⊗ phonon" space as

$$\mathcal{D}^2(\eta)=\frac{\Delta}{2\pi}\sum_\nu\left\{C_{J\nu}\left[D_j^{\lambda i}(J\nu)+F_j^{\lambda i}(J\nu)\right]\right\}^2\Big/\left[(\eta-\eta_{J\nu})^2+\Delta^2/4\right] \tag{50}$$

It is easily seen that in the limit $T\to 0$ all the Eqs. (44)-(49) transform completely into the usual formulae for strength functions in cold spherical odd nuclei obtained earlier within the QPNM framework /See, e.g., Ref./13//.

6. Schematic model

In general, to study the fragmentation of single-particle or "quasiparticle ⊗ phonon" states at finite temperature one has to solve Eq.(33) for realistic hot spherical odd nuclei. This can be done by modifying the computational procedure realized for cold nuclei in the program PHOQUS/24/ within the QPNM framework. In this paper however we do not attempt to solve this task. In order to study the qualitative effects of the quasiparticle-phonon scattering in hot odd nuclei and of the phonon scattering in hot even-even ones/11/ we employ here an oversimplified schematic model. Namely, we consider the excitation operator (see Eq.(35))

$$Q^+_{odd} = C\{\alpha^+ + D[\alpha^+ \otimes Q^+] + F[\alpha^+ \otimes \tilde{Q}]\} \qquad (51)$$

consisting of one quasiparticle α^+ with energy ε and one phonon Q^+ with energy ω for the hot odd system. Analogously, for the hot even-even system we take the excitation operator formed by two phonons with energies ω_1 and ω_2, respectively

$$Q^+_{even-even} = R\,Q^+_1 + P[Q^+_1 \otimes Q^+_2] + S[Q^+_1 \otimes \tilde{Q}_2]. \qquad (52)$$

Acting (51) and (52) on the thermal vacuum $|0,\beta\rangle_\Omega$ we obtain the excitation in hot odd and even-even systems, respectively. For simplicity, we also assume the interaction vertices to be independent of states and temperature. Each vertex therefore can be characterized by an interaction parameter Γ for the odd system and g for the even-even one. In this way we can give an upper evaluation for the phonon scattering effect at finite temperature in an even-even system because in fact the vertices W^2 due to this effect are much smaller than U^2 /See Ref./11/. The graphs corresponding to this schematic model are depicted in Fig. 3a and 3b for the even-even and odd systems, respectively. In this figure we give also the associated propagators and mass operator.

The secular equation in the given schematic model takes the form

$$E_1 - \eta - G^2\left(\frac{1 - \mathcal{N}_1 + \mathcal{N}_2}{E_1 + E_2 - \eta} + \frac{\mathcal{N}_1 + \mathcal{N}_2}{E_1 - E_2 - \eta}\right) = 0, \qquad (53)$$

where

$$E_1 = \omega_1 \;;\; E_2 = \omega_2 \;;\; \mathcal{N}_1 = -\nu_1 \;;\; \mathcal{N}_2 = \nu_2 \;;\; G = g \qquad (54)$$

for the hot even-even system, and

$$E_1 = \varepsilon \;;\; E_2 = \omega \;;\; \mathcal{N}_1 = n \;;\; \mathcal{N}_2 = \nu \;;\; G = \Gamma \qquad (55)$$

for the hot odd system.

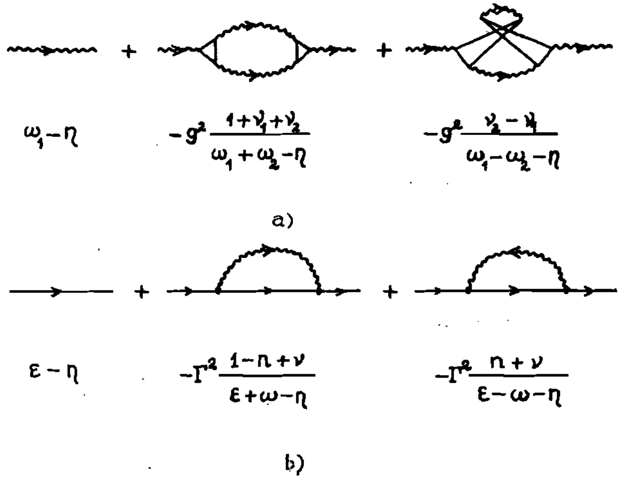

Fig. 3 Graphs corresponding to the schematic model for
a) hot even-even system
b) hot odd system.

In Eqs. (53)-(55) and in Fig. 3 the occupation numbers are

$$n = [\exp(\beta\varepsilon) + 1]^{-1}$$
$$\nu = [\exp(\beta\omega) - 1]^{-1} \qquad (56)$$
$$\nu_{1(2)} = [\exp(\beta\omega_{1(2)}) - 1]^{-1}.$$

At $T = 0$ (without scattering effect) Eq.(53) transforms into

$$E_1 - \eta - \frac{G^2}{E_1 + E_2 - \eta} = 0 \qquad (57)$$

which allows the solution

$$\eta^{(G)}_{1,2}(0) = \frac{1}{2}\left[2E_1 + E_2 \pm (E_2^2 + 4G^2)^{1/2}\right]. \qquad (58)$$

In Eq.(58) the solution $\eta^{(G)}_1(0) \equiv [2E_1 + E_2 - (E_2^2 + 4G^2)^{1/2}]/2$ is caused by the one-phonon (or single quasiparticle)component whereas the second $\eta_2^{(G)}(0)$ solution is due to the two-phonon (or "quasiparticle ⊗ phonon") term in the wave function of type (51)-(52).

When the interaction is switched off, we obtain

$$\eta_1^{(0)}(0)= E_1 \quad ; \quad \eta_2^{(0)}(0) = E_1 + E_2 . \qquad (59)$$

Therefore, we have the well-known inequalities for the repulsion of the energy levels under the influence of the interaction G:

$$\eta_1^{(G)}(0) < \eta_1^{(0)}(0) < \eta_2^{(0)}(0) < \eta_2^{(G)}(0) . \qquad (60)$$

At finite temperature $(T \neq 0)$ the solutions of Eq.(53) are defined by Cardano's formulae

$$\eta_1^{(G)}(T)=A+B+E_1 \quad ; \quad \eta_{2,3}^{(G)}(T)=-\tfrac{1}{2}(A+B)\pm i\tfrac{\sqrt{3}}{2}(A-B)+E_1 , \qquad (61)$$

where

$$A=\left[\tfrac{1}{2}G^2E_2 \tanh(\tfrac{1}{2}\beta E_1)+\sqrt{Q}\,\right]^{1/3}$$

$$B=\left[\tfrac{1}{2}G^2E_2 \tanh(\tfrac{1}{2}\beta E_1)-\sqrt{Q}\,\right]^{1/3}$$

$$Q=-\tfrac{1}{27}\left[E_2^6+3E_2^4G^2 cth(\tfrac{1}{2}\beta E_1)+G^6 cth(\tfrac{1}{2}\beta E_1)\right]- -G^4E_2^2\left[\tfrac{1}{9}cth^2(\tfrac{1}{2}\beta E_2)-\tfrac{1}{4}\tanh^2(\tfrac{1}{2}\beta E_1)\right] . \qquad (62)$$

For $Q<0$ all the three solutions (61) are real. If the interaction is absent at $T\neq 0$ we have

$$\eta_3^{(0)}(T)\equiv E_1-E_2<\eta_1^{(0)}(T)\equiv E_1<\eta_2^{(0)}(T)\equiv E_1+E_2 . \qquad (63)$$

Consequently, the inequalities analogous to (60) take the form

$$\eta_3^{(G)}(T)<\eta_3^{(0)}(T)<\eta_1^{(G)}(T)<\eta_1^{(0)}(T)<\eta_2^{(0)}(T)<\eta_2^{(G)}(T) . \qquad (64)$$

As we can see by comparing the cases with $T=0$ and $T\neq 0$, at finite temperature an additional solution $\eta_3(T)$ appears which takes the value $\eta_3^{(0)}(T)\equiv E_1-E_2$ when the interaction is absent and is shifted down with switching on the interaction. The appearance of this solution is due to the scattering process at finite temperature.

7. Numerical results

7.1. Odd system

The energies in the hot odd system calculated from Eqs. (53) and (55) for five sets of values $(E_1=\varepsilon, E_2=\omega)$ with different values of the interaction parameter Γ are displayed in Fig. 4 as functions of temperature.

From Fig. 4 it is clear that energy $\eta_1^{(\Gamma)}(T)$ converges to the value $\eta_1^{(0)}(T)\equiv\varepsilon$ of the noninteracting system with increasing T. Energy $\eta_2^{(\Gamma)}(T)$ in the set I (ω= 10 MeV, ε = 3 MeV) at first decreases with T increasing up to 4 MeV, then increases with $T > 4$ MeV. In the other set $(E_1=\varepsilon, E_2=\omega)$ energy $\eta_2^{(\Gamma)}(T)$ in general raises with increasing T. Energy $\eta_3^{(\Gamma)}(T)$ due to the quasiparticle-phonon scattering at $T\neq 0$ decreases with increasing T. The solutions obtained in the set V (ε= 10 MeV, ω = 3 MeV) turn out to be very sensitive to the variation of the interaction parameter Γ. In the set IV (ε = 7 MeV, ω = 6 MeV), for example, energy $\eta_3^{\Gamma}(T)$ goes in the negative region at $T\approx 4$ MeV while Γ = 4.0 MeV and at $T\approx 5$ MeV while Γ = 3.2 MeV.

The strength functions $C_J^2(\eta)$, $D_J^2(\eta)$ and $F_J^2(\eta)$, calculated by Eqs. (44), (45) and (46) are shown in Figs. 5 and 6, while the total strength functions $\mathcal{D}_J^2(\eta)$ (50) are dispaled in Fig. 7. From these figures, one can see that

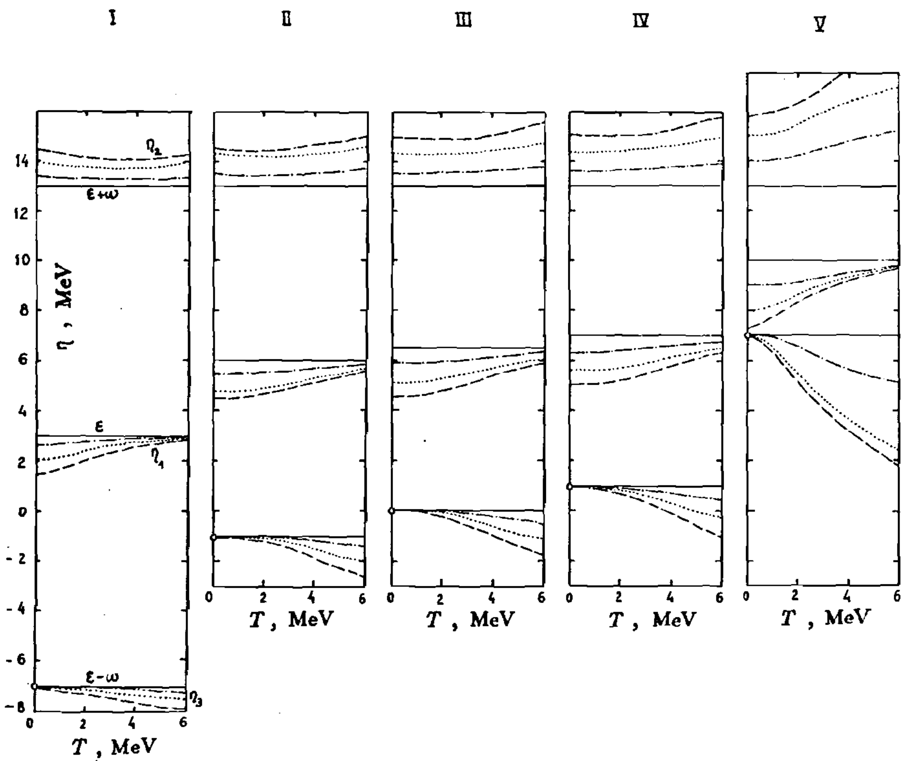

Fig. 4 Energies of the hot odd system in the schematic model calculated from Eqs. (53) and (55) for five sets (ε, ω) :

I: ε = 3 MeV, ω = 10 MeV

II: ε = 6 MeV, ω = 7 MeV

III: ε = ω = 6.5 MeV

IV: ε = 7 MeV, ω = 6 MeV

V: ε = 10 MeV, ω = 3 MeV.

The full curves denote the results obtained with $\Gamma = 0$; the dot-dashed curves are the results calculated with $\Gamma = 2$ MeV; the dotted curves stand for the results of calculations with $\Gamma = 3.2$ MeV and the dashed curves display the results when $\Gamma = 4.0$ MeV.

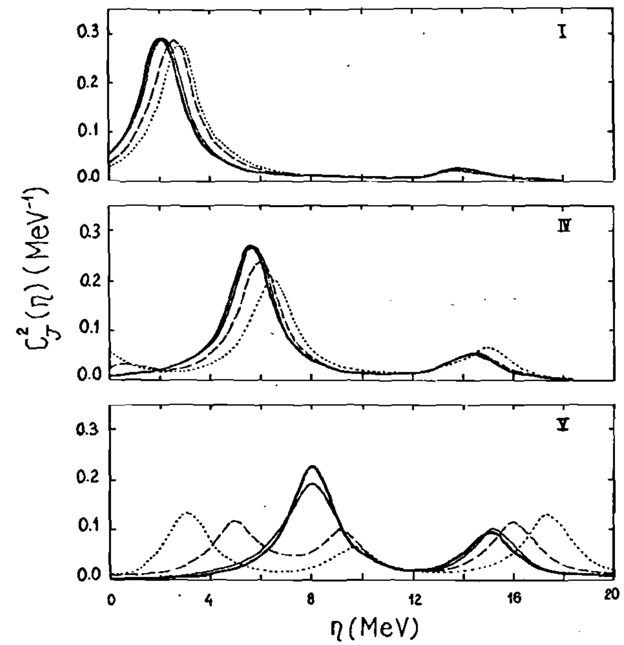

Fig. 5 Strength functions $C_J^2(\eta)$ calculated at $T = 0, 1, 3$ and 6 MeV for sets I, IV, V from the hot odd system of the schematic model in Fig. 4. The full, dashed, dotted and dot-dashed curves denote the results obtained at $T = 1, 3, 6$ and 2 MeV, respectively. The thick full curves correspond to the results at $T = 0$. $\Gamma = 3.2$ MeV.

the strength functions $C_J^2(\eta)$, $D_J^2(\eta)$ and $F_J^2(\eta)$ change not much at $T \leqslant 4 - 5$ MeV for the level $\eta_3^{(T)}(T)$ arisen by the quasiparticle-phonon scattering in the case when this level is localized near zero. In this case, the strength functions $F_J^2(\eta)$ corresponding to the quasiparticle-phonon scattering at $T \neq 0$, influence weakly the total strength functions $\mathcal{D}_J^2(\eta)$ for $T \leqslant 4 - 5$ MeV, althrough the amplitudes of $F_J^2(\eta)$ increase remarkably with T .

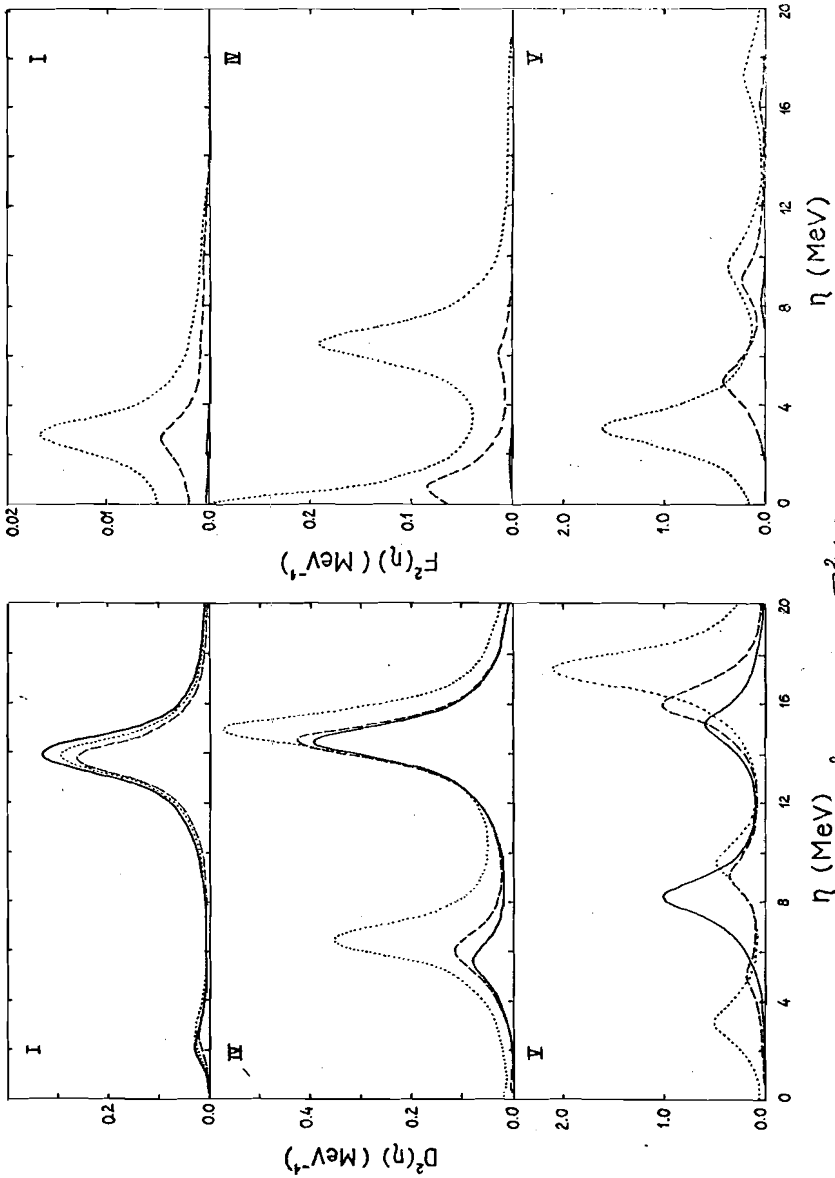

Fig.6 Strength functions $D^2_J(\eta)$ and $F^2_J(\eta)$ calculated at T=1,3 and 6 MeV for sets I,IV,V from the hot odd system of Fig. 5. The notation is as in Fig.5.

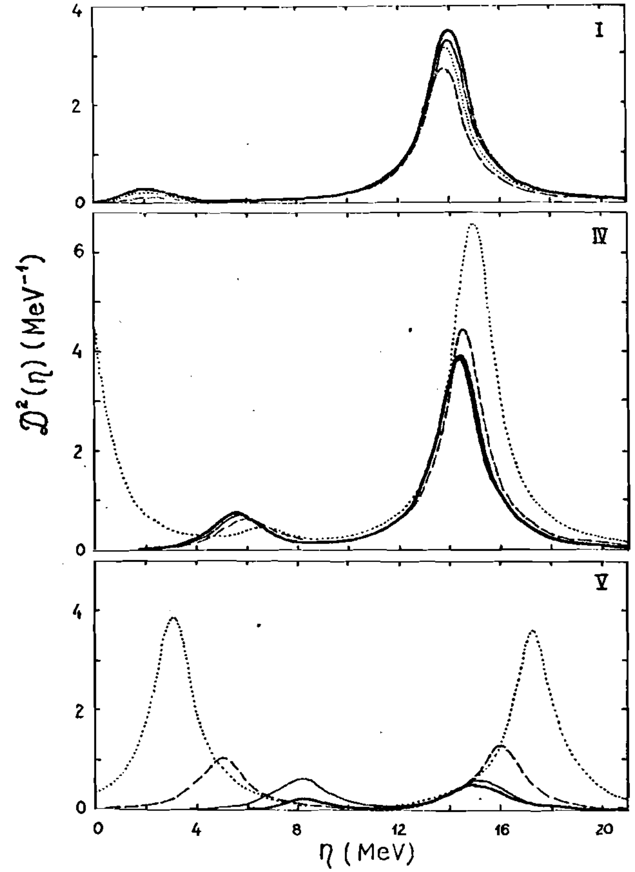

Fig.7 The strength functions $\mathcal{D}^2_J(\eta)$ calculated at $T = 0, 1, 3$ and 6 MeV (See Fig.5).

In the set V (ε = 10 MeV, ω = 3 MeV), where the value Γ = 3.2 MeV is quite strong, the peaks corresponding to $\eta^{(\Gamma)}_1(T)$ and $\eta^{(\Gamma)}_3(T)$ in the strength function $C^2_J(\eta)$ are fused in a single sizeable peak localized in 8 MeV at low temperatures (Fig.5). However, as T increases, the level $\eta^{(\Gamma)}_1(T)$ is pulsed up while $\eta^{(\Gamma)}_3(T)$ is shifted down (See fig. 4, the set V). Consequently, the broad peak at

8 MeV splits into two peaks localized at $\eta_1^{(\Gamma)}(T)$ and $\eta_3^{(\Gamma)}(T)$ going away from each other (Fig.5). We also see that the remarkable change with T is observed only for the high-lying quasiparticle levels. In the set I ($\varepsilon = 3$ MeV, $\omega = 10$ MeV) the quasiparticle level $\varepsilon = 3$ MeV is localized in the low-energy region. In this case, both the strength functions $C_J^2(\eta)$ and $D_J^2(\eta)$ change weakly with varying the temperature. For the sets IV and V, where the quasiparticle energies are large and $\varepsilon > \omega$, the strength functions $C_J^2(\eta)$, $D_J^2(\eta)$ and $\mathcal{D}_J^2(\eta)$ depend strongly upon temperature. The amplitudes of their peaks increase noticeably with increasing T up to 6 MeV. In general, with the appearance of the quasiparticle-phonon scattering in hot odd system, the single-quasiparticle and "quasiparticle $\otimes$ phonon" strength functions can change much with T for the high-lying quasiparticle levels $(\varepsilon > \omega)$.

7.2. Even-even system

In a hot even-even system one has to solve Eq. (53) having in mind Eq.(54). The solutions (61) are displayed in Fig. 8 for five sets of energies (ω_1, ω_2). The calculations have been performed with the interaction parameters g = 1.2, 2.0 and 2.8 MeV. In the even-even system the intensity the effect obtained with g = 2.8 MeV is similar to the one in the odd system with the interaction parameter Γ = 3.2 MeV. This is why we have chosen the smaller values for g in the even-even system.

Looking at Fig. 8, one can see that in contrast with the odd system, where the single-particle energy η_1 increases with T, in the even-even system the one-

phonon energy η decreasses with increasing T. The phonon scattering at finite temperature leads to the appearance of the additional level η_3 in the spectrum, obtained in the set V ($\omega_1 = 10$ MeV, $\omega_2 = 3$ MeV) In sets I, II and IV this level is negative. In set III as ω_1 is equal to ω_2, the difference in Eq. (33) vanishes and such levels do not appear($\nu_2 - \nu_1 = 0$).

In even-even nuclei, the damping of giant resonances is caused by the fragmentation of one-phonon states under the influence of the couplinge to (2p-2h) configurations. Fragmentation of that type is described by the strength function of one-phonon component. In the given schematic model this strength function has the form similar to $C_J^2(\eta)$ for the odd system where however one has to change n to $-\nu_1$, ν to ν_2, ε to ω_1, ω to ω_2 and Γ to g following Eq.(54). The strength functions of this king calculated for sets I, IV and V from Fig. 8 with the value g equal to 1.2 MeV are displayed in Fig. 9 at T =0,1,3 and 6 MeV. The noticeable change with temperature in these strength functions is observed only in set V with the high-lying one-phonon state $\omega_1 = 10$ MeV for sufficiently high values of $T \sim 6$ MeV. The value g =1,2 MeV is large enough for the two-phonon component in this set V. Moreover, the phonon scattering leads to the appearance of the additional level $\eta_3^{(g)}(T)$ associated with the peak in the region $\sim 6-7$ MeV, whose amplitude rises up with increasing T.

However as compared to the hot odd system, the phonon scattering effect in the hot even-even system is much weaker. (Compare Figs. 5 and 9 for set V).

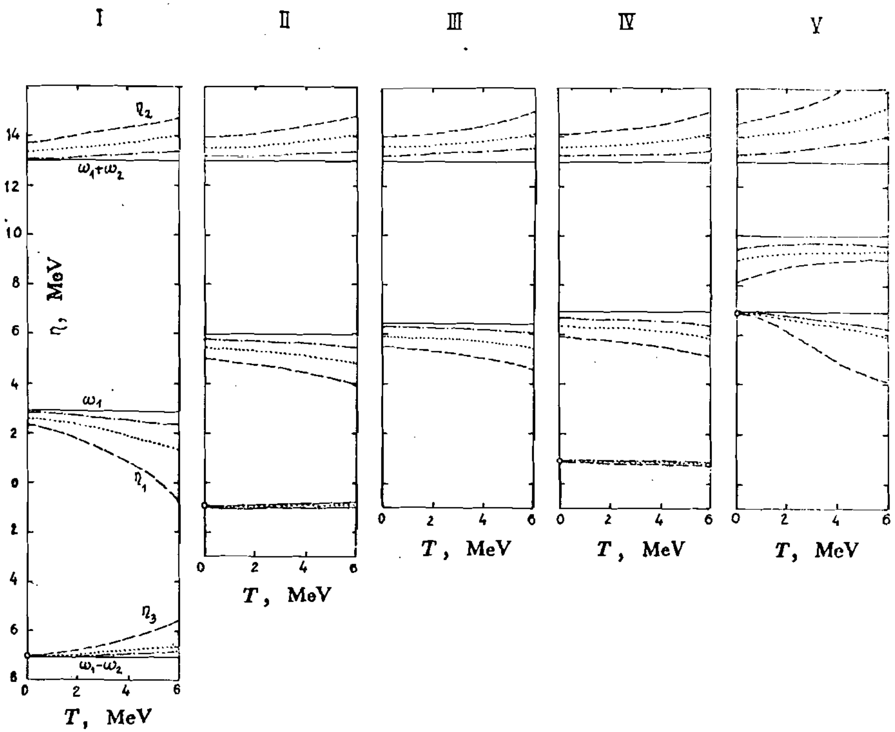

Fig. 8 Energies of the hot even-even system in the schematic model calculated from Eqs. (53) and (54) for five sets (ω_1, ω_2):

I: $\omega_1 = 3$ MeV , $\omega_2 = 10$ MeV

II: $\omega_1 = 6$ MeV , $\omega_2 = 7$ MeV

III: $\omega_1 = \omega_2 = 6.5$ MeV

IV: $\omega_1 = 7$ MeV , $\omega_2 = 6$ MeV

V: $\omega_1 = 10$ MeV , $\omega_2 = 3$ MeV .

The full curves denote the results obtained with $g = 0$; the dot-dashed curves stand for the results of calculations with $g = 1.2$ MeV and the dashed curves display the results with $g = 2.8$ MeV.

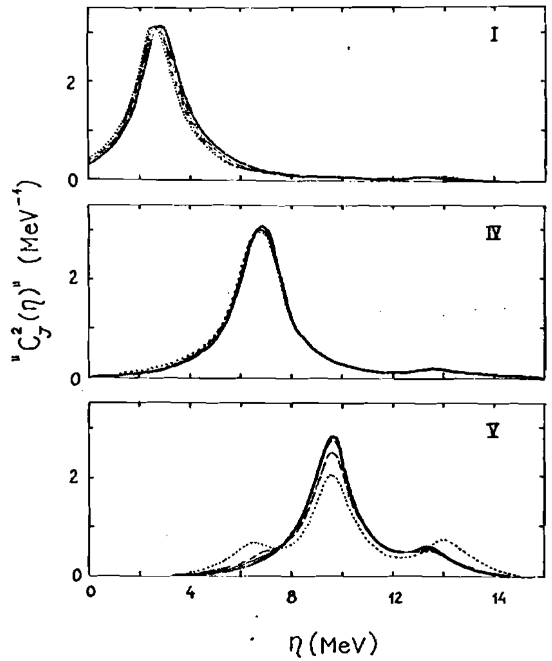

Fig. 9 Strength functions of the type $C_J^2(\eta)$ for the one-phonon component in the hot even-even system of the schematic model calculated with g =1.2 MeV at T=0,1,3 and 6 MeV for sets I,IV and V from Fig. 8. The notation is as in Fig.5.

In the calculation performed above we have used the value $U^2 = W^2 = g^2$ for the hot even-even system. In this way, we have intensified the phonon scattering effect. If we take into account the fact that $W^2 \ll U^2$, as it takes place in realistic even-even nuclei (See, Ref./11/), it is not difficult to see that the phonon scattering effect will be much weaker.

It is noteworthy also that we have also used the value for the interaction parameter Γ equal to 3.2 MeV for the hot odd system and g equal to 1.2 MeV for the even-even one. These values are in the same order with the single-quasiparticle or one-phonon energies. In practice, the vertices Γ^2 and U^2 are much weaker than the single-quasiparticle and one-phonon energies, respectively. Therefore, the scattering effects at finite temperature are expected to be much weaker as compared to the "upper limit" in the evaluation performed here. An example confirming this fact is shown in Fig. 10, where the strength function $C_J^2(\eta)$ for set V in the odd system (Fig. 4) calculated with Γ =0.4 MeV, is displayed. This set is just the one, where the quasiparticle-phonon scattering effect is the strongest. It is clear that as compared to Fig. 5, the temperature effect in Fig. 10 is negligibly small.

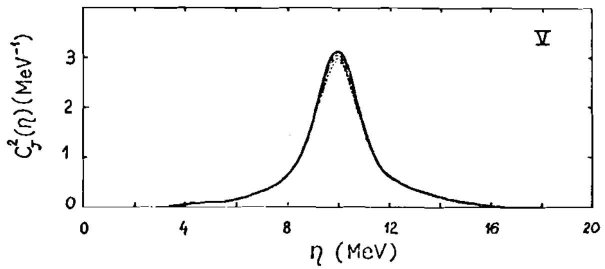

Fig. 10 The same as in Fig. 5 with Γ =0.4 MeV.

8. Conclusion

In the present work, using the temperature Green function technique we have derived the set of equations for the Green functions describing the processes with one propagating quasiparticle and the transitions from an paired quasiparticle to the coupling of intermediate phonon with quasiparticle as well as the quasiparticle-phonon scattering for hot odd spherical nuclei. The secular equation obtained within the framework of the FT-QPNM from this set defines the energies of excitations whose wave functions consist of quasiparticle and "quasiparticle $\otimes$ phonon" components. We have discussed the diagrams associated with these processes and pointed out the graphs appearing exceptionally at finite temperature.

We have also established to one-to-one correspondence between the defined Green functions and the components in the FT-QPNM excitation operators.

The numerical estimations performed in the oversimplified schematic models for hot odd and even-even systems have shown that the scattering effects at finite temperature turn out to be stronger for odd system than for even-even one. In both the types of nuclei, the phonon orquasiparticle-phonon scattering effects increase with increasing interaction and the energy of one-phonon or single-quasiparticle component. In hot even-even nuclei the phonon scattering effect at $T \neq 0$ may be noticeable only at sufficiently high temperatures beginning from $T \gtrsim$ 6 MeV. Therefore, at $T <$ 6 MeV it can be neglected in calculations. In practice, due to the small values of interaction parameters as compared to excitation energies scattering effects in hot spherical (odd and even-even) nuclei are expected to be negligibly small at moderate temperatures ($T <$ 6 MeV). The calculations based on the secular equation obtained in this work in realistic hot odd spherical nuclei will be the subject of forthcoming studies.

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