объединенный
институт
ядерных
исследований
дубна

E4-84-481

Nguyen Dinh Dang

EQUATIONS OF THE QUASIPARTICLE-PHONON NUCLEAR MODEL WITH THE PHONON SCATTERING EFFECTS AT FINITE TEMPERATURE

Submitted to "ТМФ"

1984

The recent discovery of giant resonances in heavy ion fusion[1] and deep inelastic reactions[2] has roused a growing interest for the investigation of the collective motion in highly excited nuclei. Since the transferred energy from the relative motion of colliding nuclei to the internal degrees of freedom of the reaction products undergoes rapid equilibration in times (10^{-21} sec) comparable to de-excitation times (10^{-16} sec)[3], the resulting equilibrated heavy nuclei should be described in the framework of the statistical formalism. In this approach all nuclear levels are assigned to a given temperature (an intrinsic excitation energy) with equal probabilities.

A statistical theory has been advanced in the framework of the finite temperature RPA (FT-RPA) for the description of the thermal equilibrated nuclear characteristics[4-13]. However, this theory cannot be applied to the description of the giant resonance spreading width due to the interaction with more complicated configurations in highly excited nuclei, since it is beyond the RPA-scope. In the quasiparticle-phonon nuclear model (QPNM)[14] the interaction with (2p-2h) configurations is taken into account by using the excited state wave functions, which contain two-phonon components besides one-phonon components. In this way one can describe the appearance of the giant multipole resonance (GMR) spreading width for complex cold nuclei[15]. The modification of the wave functions by including two-phonon components requires the account of the phonon-ground state correlations and the Pauli principle between phonons. Within the framework of the QPNM the phonon-ground state correlations have first been investigated for even-even spherical nuclei in Ref.[16]. In Ref.[17] a basic system of QPNM equations has been derived for even-even spherical nuclei by taking into account both the Pauli principle and phonon-ground state correlation effects at zero temperature. The diagrammatic representasion for these equations allowed one to make a comparison between the QPNM results and the results of other approaches, namely the nuclear field theory (NFT)[18-22] and the theory of finite Fermi-systems (TFFS)[23]. To obtain these equations we have assumed that the number of phonons in the ground state is small: $\langle\Psi_0|Q^+Q|\Psi_0\rangle \approx 0$[16,17]. At finite temperature ($T \neq 0$) to consider the interaction with (2p-2h) configutations we shall exactly take into account the occupation numbers of one-phonon energy levels in addition, as in the case of FT-RPA, in which the occupation numbers of quasiparticle levels are taken into account [4-13]. The structure of phonons can be calculated in the zero temperature RPA by the computing program in Ref. [24]. The occupation numbers of quasiparticle levels can be expressed in terms of the phonon operators using the well-known commutation relations [14] or the boson representation of fermion pairs [25]. The program of the present paper is the following. We are going to derive a set of equations taking into account the interaction with (2p-2h) configurations in even-even spherical nuclei at finite temperature. For simplicity the effects of the Pauli principle for two-phonon components will be approximately accounted for as in Ref. [26]. We will also analyse the new diagrams arisen at finite temperature and make a comparison with the NFT diagrams and the TFFS results.

1. DIAGRAMMATIC ANALYSIS OF ONE-PHONON PROPAGATORS AT FINITE TEMPERATURE

The QPNM Hamiltonian involves the terms describing the motion of nucleons in the average nuclear field, the superconducting pairing interaction and the residual interaction in the form of separable multipole and spin-multipole forces. In a further consideration we will study only the $E\lambda$-states generated by the multipole (p-h) forces. It is convenient to rewrite the Hamiltonian in this case in the form

$$H = \sum_{jm} \epsilon_j \alpha^+_{jm}\alpha_{jm} + H_V + H_{Vq},$$

$$H_V = -\frac{1}{4}\sum_{\lambda\mu i i' r} \frac{X_M^{\lambda i}(r) + X_M^{\lambda i'}(r)}{\sqrt{\mathcal{Y}_r^{\lambda i}\,\mathcal{Y}_r^{\lambda i'}}} \times \qquad (1)$$

$$\times\,(Q_{\lambda-\mu i} + (-)^{\lambda-\mu} Q^+_{\lambda\mu i})(Q^+_{\lambda-\mu i'} + (-)^{\lambda-\mu} Q_{\lambda\mu i'}),$$

$$H_{Vq} = -\frac{1}{2\sqrt{2}}\sum_{\lambda\mu i}\{[(-)^{\lambda-\mu}Q^+_{\lambda\mu i} + Q_{\lambda-\mu i}]\sum_{jj'r}\frac{f^{\lambda}_{jj'}\, v^{(-)}_{jj'}}{\sqrt{\mathcal{Y}_r^{\lambda i}}}\, B(jj';\lambda-\mu) + h.c.\}.$$

In Eqs. (1) ϵ_j is the quasiparticle energy; α^+_{jm} and α_{jm} are the quasiparticle creation and annihilation operators respectively; $Q^+_{\lambda\mu i}$ and $Q_{\lambda\mu i}$ are the phonon operators defined in Ref. [14]. The structure of phonons is computed in the RPA, and one can calculate the functions $X_M^{\lambda i}(r)$ and norms $\mathcal{Y}_r^{\lambda i}$ [14,24] by solving the RPA equations, where $r = \{n,p\}, -r = \{p,n\}$; $f^{\lambda}_{jj'}$ are the matrix elements for single-particle operators ge-

nerating excitations; $v^{(-)}_{jj'} = u_j u_{j'} - v_j v_{j'}$, where u_j, v_j are the Bogolubov coefficients. The operator $B(jj';\lambda-\mu)$ contains the combinations $\alpha^+\alpha$ in the form

$$B(jj';\lambda-\mu) = \sum_{mm'} (-)^{j'+m'} \langle jm\, j'm' | \lambda-\mu\rangle \alpha^+_{jm}\alpha_{j'-m'}. \qquad (2)$$

In Refs. [4-13] operator (2) has been taken into account through the occupation numbers of quasiparticle energy levels $\langle\alpha^+_{jm}\alpha_{jm}\rangle$ described by the Fermi-Dirac distribution. As we are interested in the interaction between quasiparticles and phonons described by the terms H_{Vq} of Eqs. (1), it is convenient to use the representation, in which the occupation numbers of one-phonon levels are taken exactly into account. This is carried out as follows. Using the commutation relations between operators B and Q^+ [14], we expand the Hamiltonian (1) over phonon operators Q, Q^+, as has been done in Ref. [16]

$$H = \sum_{\lambda\mu i}\omega_{\lambda i} Q^+_{\lambda\mu i}Q_{\lambda\mu i} + \frac{1}{2}\sum_{\substack{\lambda\mu i\\ \lambda_1\mu_1 i_1\\ \lambda_2\mu_2 i_2}} \langle\lambda_1\mu_1\lambda_2\mu_2|\lambda-\mu\rangle \times \qquad (3)$$

$$\times\,[\,U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)\, Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2} Q_{\lambda-\mu i} +$$

$$+ (-)^{\lambda-\mu} V^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)\, Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2} Q^+_{\lambda\mu i} + h.c.].$$

Eq. (3) can be rewritten in the form with normal and anomalous components for the terms containing $U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)$:

$$H = \sum_{\lambda\mu i}\omega_{\lambda i} Q^+_{\lambda\mu i}Q_{\lambda\mu i} + \frac{1}{2}\sum_{\substack{\lambda\mu i\\ \lambda_1\mu_1 i_1\\ \lambda_2\mu_2 i_2}} \langle\lambda_1\mu_1\lambda_2\mu_2|\lambda-\mu\rangle \times \qquad (4)$$

$$\times\,[\,\tilde{U}^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i) Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2} Q_{\lambda-\mu i} + (-)^{\lambda_1-\mu_1} W^{\lambda i}_{\lambda_1 i_1}(\lambda_2 i_2)\, Q^+_{\lambda_1\mu_1 i_1} Q_{\lambda_2-\mu_2 i_2} Q^+_{\lambda\mu i} +$$

$$+ (-)^{\lambda-\mu} V^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)\, Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2} Q^+_{\lambda\mu i} + h.c.].$$

Expressions for the coefficients $U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)$ and $V^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)$ from Eqs. (3) and (4) are given in Refs. [16,17]. Coefficients

$\bar U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)$ and $W^{\lambda i}_{\lambda_1 i_1}(\lambda_2 i_2)$ are defined as follows:

$$\bar U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i) = (-)^{\lambda_1+\lambda_2-\lambda}\sqrt{(2\lambda_1+1)(2\lambda_2+1)}\ \times$$
$$\times \sum_{j_1 j_2 j_3}\Big[\frac{f^{\lambda_1}_{j_1 j_2}\, v^{(-)}_{j_1 j_2}}{\sqrt{2\mathcal{Y}_{\lambda_1 i_1}}}\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda\\ j_3 & j_2 & j_1\end{Bmatrix}(\psi^{\lambda i}_{j_2 j_3}\psi^{\lambda_2 i_2}_{j_3 j_1} + \phi^{\lambda i}_{j_2 j_3}\phi^{\lambda_2 i_2}_{j_3 j_1}) +$$
$$+ \frac{f^{\lambda_2}_{j_1 j_2}\, v^{(-)}_{j_1 j_2}}{\sqrt{2\mathcal{Y}_{\lambda_2 i_2}}}\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda\\ j_1 & j_3 & j_2\end{Bmatrix}(\psi^{\lambda i}_{j_3 j_1}\psi^{\lambda_1 i_1}_{j_2 j_3} + \phi^{\lambda i}_{j_3 j_1}\phi^{\lambda_1 i_1}_{j_2 j_3})\Big], \qquad (5)$$

$$W^{\lambda i}_{\lambda_1 i_1}(\lambda_2 i_2) = (-)^{\lambda_1+\lambda_2-\lambda}\sqrt{(2\lambda_1+1)(2\lambda_2+1)}\ \times$$
$$\times \sum_{j_1 j_2 j_3}\frac{f^{\lambda_2}_{j_1 j_2}\, v^{(-)}_{j_1 j_2}}{\sqrt{2\mathcal{Y}_{\lambda_2 i_2}}}\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda\\ j_2 & j_3 & j_1\end{Bmatrix}(\psi^{\lambda_1 i_1}_{j_3 j_1}\phi^{\lambda i}_{j_2 j_3} + \phi^{\lambda_1 i_1}_{j_3 j_1}\psi^{\lambda i}_{j_2 j_3}).$$

The relation of coefficients (5) with $U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)$ is straightforward

$$\bar U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)\langle\lambda_1\mu_1\lambda_2\mu_2|\lambda-\mu\rangle + (-)^{\lambda_1-\mu_1}W^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)\langle\lambda_1\mu_1\lambda\mu|\lambda_2-\mu_2\rangle =$$
$$= U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda i)\langle\lambda_1\mu_1\lambda_2\mu_2|\lambda-\mu\rangle. \qquad (6)$$

The diagrammatic representation for the terms with U, Ũ, W and V from Eq. (4) is given by Fig. 1. In Refs. [16,17] an approximation neglecting the number of phonons in the ground state has been used, which corresponds to the zero temperature limit. By analogy, in Ref. [27] the zero temperature Green function method has been employed to one-phonon propagators. In both cases these approximations are equivalent to the disregard of the anomalous components with W from Eq. (4), and therefore, only the vertices U and V from Fig. 1 were taken into account. At finite temperature the occupation numbers of one-phonon levels depend on the one-phonon energies $\omega_{\lambda i}$ and the temperature T according to the Bose distribution

$$\nu_{\lambda i} \equiv \langle Q^+_{\lambda\mu i} Q_{\lambda\mu i}\rangle = [\exp(\omega_{\lambda i}/T) - 1]^{-1}. \qquad (7)$$

In this equation the angular brackets $\langle\dots\rangle$ denote the thermal averages in the grand canonical ensemble [28,29]. The energies

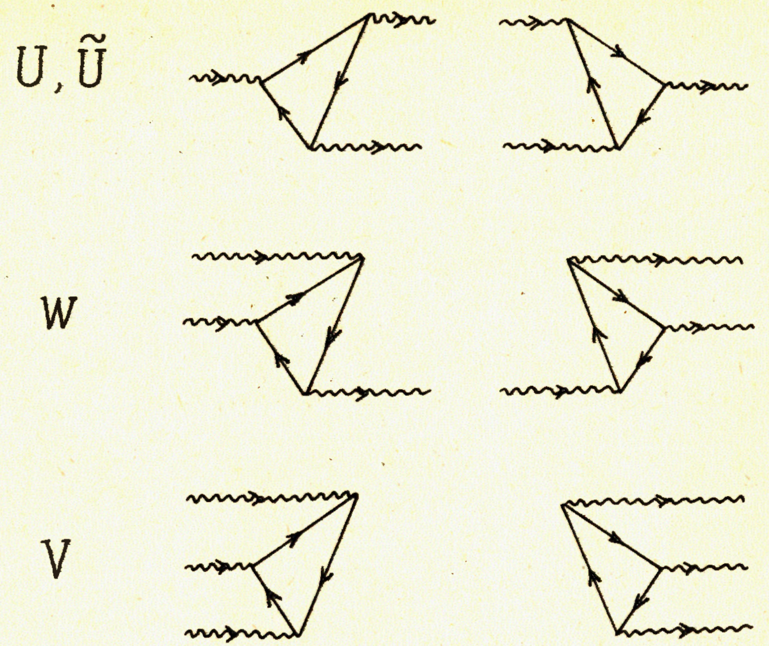

Fig. 1. Vertices **U, Ũ, W** *and* **V**. *Wavy lines are phonons; solid lines are quasiparticles.*

$\omega_{\lambda i}$ of phonons are calculated in the RPA [14,24] and the temperature T is given in MeV.

According to the formalism of finite temperature Green functions advanced in Refs. [28,29], one can define the unperturbed Green function with the time depended phonon operators $Q^+_{\lambda\mu i}(t)$ and $Q_{\lambda\mu i}(t)$ [27,29]

$$G^{(0)}_{\lambda\mu i,\lambda'\mu' i'}(t-t') = -i\langle T\{Q_{\lambda\mu i}(t)\,Q^+_{\lambda'\mu' i'}(t')\}\rangle. \qquad (8)$$

Relating $Q_{\lambda\mu i}(t)$ and $Q^+_{\lambda'\mu' i'}(t)$ with the time independent phonon operators in the Schrödinger picture and consulting Eq. (7), we use the generalized Wick theorem for canonical products [30] to obtain

$$G^{(0)}_{\lambda\mu i,\lambda'\mu' i'}(t-t') = -i\{\theta(t-t')\exp[-i\omega_{\lambda i}(t-t')](1+\nu_{\lambda i}) +$$
$$+ \theta(t'-t)\exp[i\omega_{\lambda i}(t'-t)]\nu_{\lambda i}\}\delta_{\lambda\lambda'}\delta_{\mu\mu'}\delta_{ii'}. \qquad (9)$$

Function (9) has the Fourier component

$$G^{(0)}_{\lambda\mu i,\lambda'\mu' i'}(\eta) = \Big[\frac{1+\nu_{\lambda i}}{\eta-\omega_{\lambda i}+i\epsilon} - \frac{\nu_{\lambda i}}{\eta+\omega_{\lambda i}-i\epsilon}\Big]\delta_{\lambda\lambda'}\delta_{\mu\mu'}\delta_{ii'}. \qquad (10)$$

From Eqs.(9) and (10) one can easily see that the unperturbed propagator $G^{(0)}$ used in Ref. [27] is the particular case of functions (9) and (10) in the zero temperature limit, in which the canonical products are reduced to the normal one (Cf.Ref. [30]). Similar functions have been applied in TFFS [23] to construct two-phonon diagrams involving the interaction with (2p - 2h) configurations.

By consulting Eqs.(9), (10) we now expand in the interaction picture the exact Green function in a perturbation theory series as in Ref. [29]

$$G_{JMi,JMi'}(t-t') = G^{(0)}_{ii'}(JM;t-t') + \sum_{n=1}^{\infty} \delta G^{(2n)}_{ii'}(JM;t-t') =$$

$$= G^{(0)}_{ii'}(JM;t-t') + (-i)\sum_{n=1}^{\infty} \frac{(-i)^{2n}}{2n!} \int_{-\infty}^{+\infty} \langle T\{Q_{JMi}(t) \times \tag{11}$$

$$\times Q^{+}_{JMi'}(t') H_{Vq}(r_1) \dots H_{Vq}(r_{2n})\}\rangle dr_1 \dots dr_{2n},$$

and employing the above-mentioned generalized Wick theorem for canonical products, we find for the Fourier components $\delta G^{(2)}_{ii'}(JM;\eta)$

$$\delta G^{(2)}_{ii'}(JM;\eta) = [\hat{G}^{(0)}(JM;\eta)\,\hat{\Pi}^{U+V+W}(J;\eta)\,\hat{G}^{(0)}(JM;\eta)]_{ii'}. \tag{12}$$

Eq.(12) provides a generalized presentation of the functions $\delta G^{(2)}_{ii'}(JM;\eta)$, for which the propagators in the right-hand side have the form

$$\hat{G}^{(0)}_{ii'}(JM;\eta) = [\frac{1+\nu_{Ji}}{\eta-\omega_{Ji}+i\epsilon} - \frac{\nu_{Ji}}{\eta+\omega_{Ji}-i\epsilon}]\delta_{ii'} \equiv \hat{G}^{(0)}_{ii'}(J;\eta),$$

$$\hat{\Pi}^{U+V+W}_{ii'}(J;\eta) = \hat{\Pi}^{U}_{ii'}(J;\eta) + \hat{\Pi}^{V}_{ii'}(J;\eta) + \hat{\Pi}^{W}_{ii'}(J;\eta).$$

$$\hat{\Pi}^{U}_{ii'}(J;\eta) = \sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} \frac{U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)\,U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji')}{\eta-\omega_{\lambda_1 i_1}-\omega_{\lambda_2 i_2}}(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2}), \tag{13}$$

$$\hat{\Pi}^{V}_{ii'}(J;\eta) = \sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} \frac{V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)\,V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji')}{\eta+\omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2}}(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2}),$$

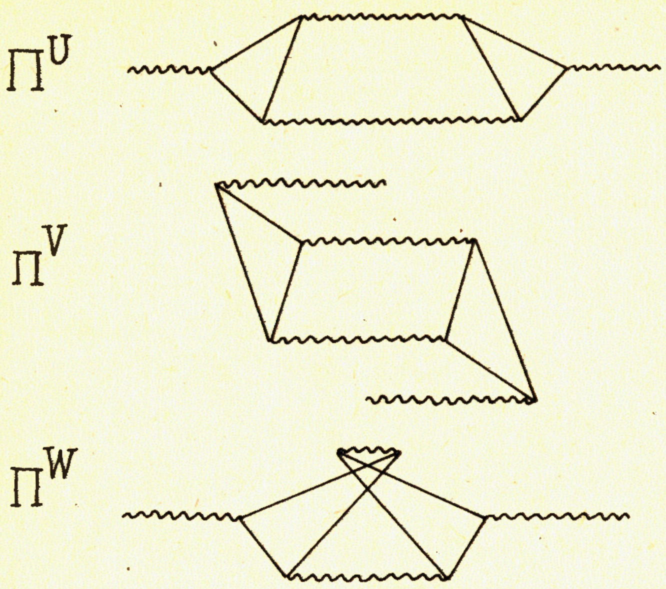

Fig.2. The proper self-energies Π^U, Π^V and Π^W from Eqs.(13).

$$\hat{\Pi}^{W}_{ii'}(J;\eta) = \sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} [\frac{W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)\,W^{Ji'}_{\lambda_1 i_1}(\lambda_2 i_2)}{\eta+(\omega_{\lambda_1 i_1}-\omega_{\lambda_2 i_2})} - \frac{W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)\,W^{Ji'}_{\lambda_2 i_2}(\lambda_1 i_1)}{\eta-(\omega_{\lambda_1 i_1}-\omega_{\lambda_2 i_2})}] \times$$

$$\times(\nu_{\lambda_2 i_2}-\nu_{\lambda_1 i_1}).$$

In the zero temperature limit $T \to 0$ one can derive immediately from Eqs.(13) the proper self-energy $\hat{\Pi}^{U+V+W}_{ii'}(J;\eta) \to \hat{\Pi}^{U+V}_{ii'}(J;\eta)$ obtained in Ref. [27].

Following the diagrammatic correspondence given in Ref. [17] and the vertices in Fig.1 we can illustrate the proper self-energies Π^U, Π^V and Π^W from Eqs.(13) by means of the diagrammatic representation given in Fig.2.

Let us now analyse the consequences caused by the diagrams Π^W which appear only at finite temperature when the occupation numbers $\nu_{\lambda i}$ of one-phonon levels are taken exactly into account. Like the appearance of the new particle-particle (p - p) and hole-hole (h - h) poles in the FT-RPA discussed in Refs. [4,5,8], the new terms Π^W in the self-energy operator lead to the appearance of poles of the type $\pm(\omega_{\lambda_1 i_1}-\omega_{\lambda_2 i_2})$ besides the two-phonon poles at energies $(\omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2})$ originating from Π^U. The form of these proper self-energies enables us to identify the vertices U, V and W as the processes of one- to two-phonon transitions and vice versa, the phonon correlations and the phonon scattering, respectively. We also note that similar terms have been obtained in the framework of the TFFS [23] by a somewhat different method.

Table

The $\Pi^U_{ii'}$-, $\Pi^V_{ii'}$- and $\Pi^W_{ii'}$-values (in MeV) from the proper self-energy $\Pi^{U+V+W}_{ii'}$ of the Dyson equation $G = G^{(0)} + G^{(0)}\Pi^{U+V+W}G$ for excitations built on collective states 2^+_1 with energy $\omega_{2^+_1} = 1.6554$ MeV and 2^+_6 with energy $\omega_{2^+_6} = 3.6183$ MeV in the ^{118}Sn nucleus. SUM = $|\Pi^U_{ii'}| + |\Pi^V_{ii'}| + |\Pi^W_{ii'}|$

$\Pi^U_{ii'}$	$\Pi^V_{ii'}$	$\Pi^W_{ii'}$
0,1723	$0,4741.10^{-2}$	$-0,1013.10^{-1}$
$-0,1262.10^{-1}$	$0,2501.10^{-3}$	$0,1927.10^{-2}$
$-0,1262.10^{-1}$	$0,2501.10^{-3}$	$0,1927.10^{-2}$
$0,7539.10^{-1}$	$0,1494.10^{-4}$	$-0,2434.10^{-3}$

SUM	$\|\Pi^U_{ii'}\|$/SUM	$\|\Pi^V_{ii'}\|$/SUM	$\|\Pi^W_{ii'}\|$/SUM
0,1871	92%	2,5%	5,4%
0,1480	85%	1,7%	1,3%
0,1480	85%	1,7%	1,3%
0,7565	99,7%	<1%	<1%

To evaluate the contribution of the new roots between the poles at energies $\pm(\omega_{\lambda_1 i_1} - \omega_{\lambda_2 i_2})$ we note that due to the exponential dependence of the numerators in Π^W(13) on the values $\omega_{\lambda i}/T$, the values of the reduced probabilities $B(E\lambda)$ for the electric transitions associated with these levels turn out to be negligible as compared with the $B(E\lambda)$-values for the (p-h) levels. To illustrate this statement we show the result obtained by calculating a schematic case for collective modes built on

two collective 2^+_1 and 2^+_6-states with energies $\omega_{2^+_1} = 1.655$ MeV and $\omega_{2^+_6} = 3.618$ MeV in the ^{118}Sn nucleus. A more complete information about the classification of collective states in the QPNM can be found in Ref.[15]. In this evaluation we choose the upper boundary for the function Π^W putting $\nu_{\lambda_2 i_2} - \nu_{\lambda_1 i_1} \approx 1$. In general the terms $\Pi^W_{ii'}$ are functions of excitation energies η, but it is enough to calculate the values of $\Pi^W_{ii'}$ at a fixed value of η. The calculated results for $\eta = 1$ MeV are collected in the Table. One can conclude that the effects associated with the phonon-ground state correlations and the phonon scattering at finite temperature are weak comparable with the one- to two-phonon transitions.

2. SYSTEM OF APPROXIMATE EQUATIONS FOR GREEN FUNCTIONS

We are going now to derive a generalized set of equations for Green functions taking exactly into account the occupation numbers of one-phonon levels (7). To realize this idea besides the Green functions employed in Ref.[27]

$$\begin{aligned}
G^{-,+}_{ii'}(JM; t-t') &= \langle\langle Q_{JMi}(t), Q^+_{JMi'}(t')\rangle\rangle, \\
G^{+,+}_{ii'}(JM; t-t') &= \langle\langle Q^+_{JMi}(t), Q^+_{JMi'}(t')\rangle\rangle, \\
G^{--,+}_{12,3}(t-t') &= \langle\langle Q_{\lambda_1\mu_1 i_1}(t)\, Q_{\lambda_2\mu_2 i_2}(t), Q^+_{\lambda_3\mu_3 i_3}(t')\rangle\rangle, \\
G^{++,+}_{12,3}(t-t') &= \langle\langle Q^+_{\lambda_1\mu_1 i_1}(t)\, Q^+_{\lambda_2\mu_2 i_2}(t), Q^+_{\lambda_3\mu_3 i_3}(t')\rangle\rangle
\end{aligned} \qquad (14)$$

we have to consider two anomalous Green functions

$$\begin{aligned}
G^{+-,+}_{12,3}(t-t') &= \langle\langle Q^+_{\lambda_1\mu_1 i_1}(t)\, Q_{\lambda_2\mu_2 i_2}(t), Q^+_{\lambda_3\mu_3 i_3}(t')\rangle\rangle, \\
G^{-+,+}_{12,3}(t-t') &= \langle\langle Q_{\lambda_1\mu_1 i_1}(t)\, Q^+_{\lambda_2\mu_2 i_2}(t), Q^+_{\lambda_3\mu_3 i_3}(t')\rangle\rangle = G^{+-,+}_{21,3}(t-t')
\end{aligned} \qquad (15)$$

which naturally appear at finite temperature. The double angular brackets in the right-hand sides of Eqs.(14) and (15) at finite temperature correspond to the thermal averages of canonical products.

By the usual method for Green functions[28] we obtained a set of coupled equations for the Fourier components of functions (14), (15):

$$(\omega_{Ji}-\eta)G^{-,+}_{ii'}(JM;\eta)+\frac{1}{2}\sum_{\substack{\lambda_1\mu_1 i_1\\ \lambda_2\mu_2 i_2}}\{\langle\lambda_1\mu_1\lambda_2\mu_2|JM\rangle[U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)G^{--,+}_{12,JMi'}(\eta)+$$
$$+3(-)^{J-M}V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)G^{++,+}_{12,JMi'}(\eta)]+\langle\lambda_1\mu_1\lambda_2-\mu_2|JM\rangle[(-)^{\lambda_1-\mu_1}W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)G^{+-,+}_{12,JMi'}(\eta)+$$
$$+(-)^{\lambda_2-\mu_2}(-)^{J-M}W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)G^{+-,+}_{21,JMi'}(\eta)]\}=-\delta_{ii'},$$

$$(\omega_{Ji}+\eta)G^{+,+}_{ii'}(JM;\eta)+\frac{1}{2}\sum_{\substack{\lambda_1\mu_1 i_1\\ \lambda_2\mu_2 i_2}}\{\langle\lambda_1\mu_1\lambda_2\mu_2|JM\rangle[U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)G^{++,+}_{12,JMi'}(\eta)+$$
$$+3(-)^{J-M}V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)G^{--,+}_{12,JMi'}(\eta)+\langle\lambda_1\mu_1\lambda_2-\mu_2|JM\rangle[(-)^{\lambda_1-\mu_1}W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)G^{+-,+}_{21,JMi'}(\eta)+$$
$$+(-)^{\lambda_2-\mu_2}(-)^{J-M}W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)G^{+-,+}_{12,JMi'}(\eta)]\}=0, \qquad (16)$$

$$(\omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2}-\eta)G^{--,+}_{12,JMi'}(\eta)+\sum_i\langle\lambda_1\mu_1\lambda_2\mu_2|JM\rangle(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2})\times$$
$$\times[U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)G^{-,+}_{ii'}(JM;\eta)+\frac{3}{2}(-)^{J-M}V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)G^{+,+}_{ii'}(JM;\eta)]=0,$$

$$(\omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2}+\eta)G^{++,+}_{12,JMi'}(\eta)+\sum_i\langle\lambda_1\mu_1\lambda_2\mu_2|JM\rangle(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2})\times$$
$$\times[U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)G^{+,+}_{ii'}(JM;\eta)+\frac{3}{2}(-)^{J-M}V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)G^{-,+}_{ii'}(JM;\eta)]=0,$$

$$(\omega_{\lambda_2 i_2}-\omega_{\lambda_1 i_1}-\eta)G^{+-,+}_{12,JMi'}(\eta)+\sum_i(\nu_{\lambda_1 i_1}-\nu_{\lambda_2 i_2})[(-)^{\lambda_2-\mu_2}\langle\lambda_2\mu_2\lambda_1-\mu_1|JM\rangle\times$$
$$\times W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)G^{+,+}_{ii'}(JM;\eta)+(-)^{\lambda_1-\mu_1}\langle\lambda_1\mu_1\lambda_2-\mu_2|JM\rangle W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)G^{-,+}_{ii'}(JM;\eta)]=0,$$

$$(\omega_{\lambda_2 i_2}-\omega_{\lambda_1 i_1}+\eta)G^{+-,+}_{21,JMi'}(\eta)+\sum_i(\nu_{\lambda_1 i_1}-\nu_{\lambda_2 i_2})[(-)^{\lambda_1-\mu_1}\langle\lambda_1\mu_1\lambda_2-\mu_2|JM\rangle\times$$
$$\times W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)G^{+,+}_{ii'}(JM;\eta)+(-)^{\lambda_2-\mu_2}\langle\lambda_2\mu_2\lambda_1-\mu_1|JM\rangle W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)G^{-,+}_{ii'}(JM;\eta)]=0.$$

In the framework of the Hamiltonians (3) and (4) with the condition for closing the chain of Green functions at two-phonon terms (14) and (15), system (16) is exact. The Pauli principle between two-phonon components can be taken here into account exactly in calculations using the exact phonon commutation relations as in Refs./15,17/. In the zero temperature limit the

numbers $\nu_{\lambda_1 i_1}$ and $\nu_{\lambda_2 i_2}$ in Eqs.(16) become zero, so the two last equations from Eqs.(16) are identities and the other four form the system obtained ealier in Ref. /27/. The Green functions (14) and (15) can be put in the one-to-one correspondence with the coefficients of the QPNM excited state wave function:

$$G^{-,+}_{ii'}\leftrightarrow\xi^J_{i'}R^{i'}_i(J), \qquad G^{++,+}_{12,JMi}\leftrightarrow 2\zeta^J_{i'}P^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji),$$
$$G^{+,+}_{ii'}\leftrightarrow\zeta^J_{i'}R^{i'}_i(J), \qquad G^{+-,+}_{12,JMi}\leftrightarrow 2\xi^J_{i'}S^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji), \qquad (17)$$
$$G^{--,+}_{12,JMi}\leftrightarrow 2\xi^J_{i'}P^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji), \qquad G^{+-,+}_{21,JMi}\leftrightarrow 2\zeta^J_{i'}S^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji).$$

Accordingly, the excitation operator can be defined in the form

$$\mathfrak{E}^+_I(JM)=\sum_\nu[\xi^{JI}_\nu(\Omega^+_{JM\nu}+\Delta^+_{JM\nu})-(-)^{J-M}\zeta^{JI}_\nu(\Omega_{J-M\nu}+\Delta_{J-M\nu})], \qquad (18)$$

where

$$\Omega^+_{JM\nu}=\sum_i R^\nu_i(J)Q^+_{JMi}+\sum_{\substack{\lambda_1\mu_1 i_1\\ \lambda_2\mu_2 i_2}}\langle\lambda_1\mu_1\lambda_2\mu_2|JM\rangle P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)Q^+_{\lambda_1\mu_1 i_1}Q^+_{\lambda_2\mu_2 i_2},$$
$$\Delta^+_{JM\nu}=\sum_{\substack{\lambda_1\mu_1 i_1\\ \lambda_2\mu_2 i_2}}(-)^{\lambda_1-\mu_1}\langle\lambda_1\mu_1\lambda_2-\mu_2|JM\rangle S^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)Q^+_{\lambda_1\mu_1 i_1}Q_{\lambda_2\mu_2 i_2}. \qquad (19)$$

The energies η_{JI} of excited states can be found by linearizing the following equations of motion:

$$[H,\mathfrak{E}^+_I(JM)]=\eta_{JI}\mathfrak{E}^+_I(JM), \qquad [H,\mathfrak{E}_I(JM)]=-\eta_{JI}\mathfrak{E}_I(JM). \qquad (20)$$

As in Refs./16,17/ we obtain at first the set of equations for finding the coefficients ξ, ζ, R, P and S. This set is equivalent to system (16):

$$\sum_\nu\{\xi^{JI}_\nu[(\omega_{Ji}-\eta)R^\nu_i(J)+\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}}U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)+\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}}W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)S^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)]$$
$$+\zeta^{JI}_\nu[3\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}}V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)+\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}}W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)S^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)]\}=0,$$

$$
\begin{aligned}
&\sum_{\nu}\{\zeta^{JI}_{\nu}[(\omega_{Ji}+\eta)R^{\nu}_{i}(J)+\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)\,P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)+\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)S^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)]+\\
&+\xi^{JI}_{\nu}[3\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)+\sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)S^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)]\}=0\,,\\
&\sum_{\nu}\{\xi^{JI}_{\nu}[(\omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2}-\eta)P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)+\frac{1}{2}\sum_{i} U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)\,(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2})R^{\nu}_{i}(J)]+\\
&+\frac{3}{2}\zeta^{JI}_{\nu}\sum_{i} V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)\,(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2})R^{\nu}_{i}(J)\}=0,\\
&\sum_{\nu}\{\zeta^{JI}_{\nu}[(\omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2}+\eta)P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)+\frac{1}{2}\sum_{i} U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2})R^{\nu}_{i}(J)]+ \qquad (21)\\
&+\frac{3}{2}\xi^{JI}_{\nu}\sum_{i} V^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)(1+\nu_{\lambda_1 i_1}+\nu_{\lambda_2 i_2})R^{\nu}_{i}(J)\}=0\,,\\
&\sum_{\nu}\{\xi^{JI}_{\nu}[(\omega_{\lambda_2 i_2}-\omega_{\lambda_1 i_1}-\eta)S^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)+\frac{1}{2}\sum_{i} W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)(\nu_{\lambda_1 i_1}-\nu_{\lambda_2 i_2})R^{\nu}_{i}(J)]+\\
&+\frac{1}{2}\zeta^{JI}_{\nu}\sum_{i} W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)(\nu_{\lambda_1 i_1}-\nu_{\lambda_2 i_2})R^{\nu}_{i}(J)\}=0\,,\\
&\sum_{\nu}\{\zeta^{JI}_{\nu}[(\omega_{\lambda_2 i_2}-\omega_{\lambda_1 i_1}+\eta)S^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)+\frac{1}{2}\sum_{i} W^{Ji}_{\lambda_1 i_1}(\lambda_2 i_2)(\nu_{\lambda_1 i_1}-\nu_{\lambda_2 i_2})R^{\nu}_{i}(J)]+\\
&+\frac{1}{2}\xi^{JI}_{\nu}\sum_{i} W^{Ji}_{\lambda_2 i_2}(\lambda_1 i_1)(\nu_{\lambda_1 i_1}-\nu_{\lambda_2 i_2})R^{\nu}_{i}(J)\}=0,
\end{aligned}
$$

$i = 1, 2, \ldots, n_1$; $\nu, I = 1, 2, \ldots, (n_1+n_2)$, where n_1 and n_2 are, respectively, the numbers of one- and two-phonon components of the wave function (18).

The calculation of energies η from Eqs.(21) is straightforward by solving the secular equations for the determinant of the left-hand sides of Eqs.(21) $\det\|M\| = 0$, where M is the matrix for the set of Eqs.(21), At zero temperature system (21) is completely transformed into the system of equations derived in Ref./16/.

Systems (16) and (21) correspond to the couplings represented by the diagrams in Fig.3a. In the transition from the QPNM to the NFT diagrams the intermediate noncollective phonons are replaced by two quasiparticles (Cf. Refs. /15,17/) and the proper self-energy Π^W is transformed into the diagrams on the right from arrow in Fig.3b. At the same time, the proper self-energies

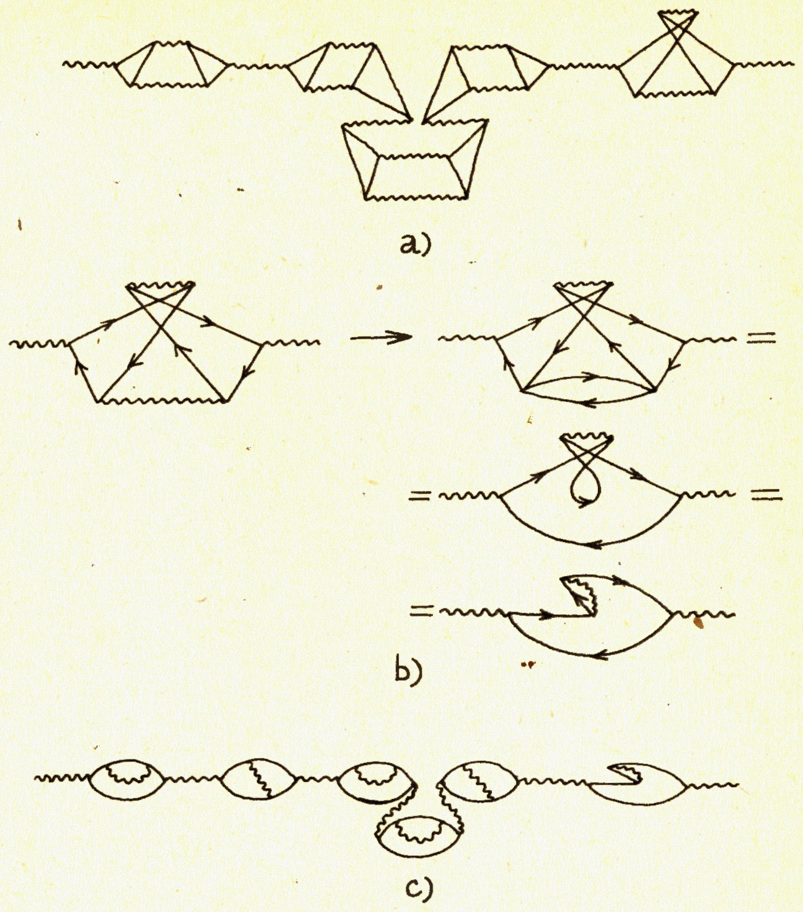

Fig.3. The QPNM and NFT diagrams. a - The QPNM diagrammatic chain. b - Rewriting the Π^W diagram in the QPNM (left from arrow) into the couplings used in the NFT (right from arrow). c - The NFT diagrammatic chain.

Π^U and Π^V are changed as in Refs./15,17/ and we obtain the resulting chain of summarized NFT diagrams /22/ given by Fig.3c. In these diagrams, as has been shown in Ref. /17/, the diagrams representing the one- to two-phonon transitions and vice versa in the QPNM are the main ones(Fig.4a). In the transition to the NFT we obtain from these diagrams the couplings with one intermediate phonon (Fig.4b).

Diagrams of the type of Fig.3b in the NFT appear because of mixing in the excited states, besides the components with one intermediate phonon, components consisting of two particle-

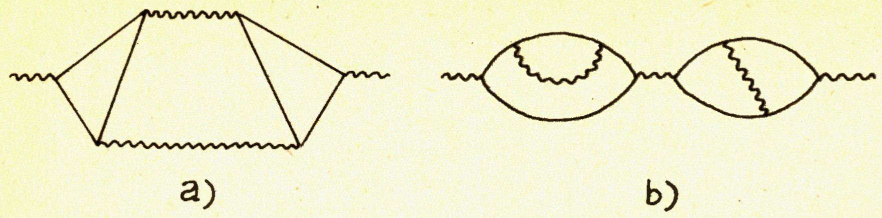

Fig.4. Main diagrams used in the QPNM (a) and in the NFT (b).

hole pairs. Such interactions cannot be interpreted as interactions between (p-h) pairs /31/. It has been shown in Ref./8/ that interactions of this kind are conditioned also by the couplings of (p-p) and (h-h) types which are caused in their turn by taking exactly into account the occupation numbers of quasiparticle energy levels at finite temperature. Thereby, neither the (p-p) and (h-h) poles in the case of simple (1p-1h) configurations nor the poles appeared at energies $\pm(\omega_{\lambda_1 i_1} - \omega_{\lambda_2 i_2})$ in the case of interaction with (2p-2h) configurations can be correctly investigated without the introduction of the temperature for highly excited nuclei as a physical quantity which can be defined from experiments. However, one can make also an evaluation of the occupation numbers of quasiparticle levels in the perturbation theory at zero temperature following the introduction of phonons as in Ref. /14/. In this case these numbers are thought to be those of quasiparticles in the phonon-ground state. The evaluation of this kind has been done in Ref. /32/ for deformed nuclei, and its results have shown that the numbers of quasiparticles in the phonon-ground state are negligible. For even-even spherical nuclei the contribution of terms containing B^+B to the excitation energies has been calculated in the perturbation theory for the RPA in Ref. /33/ and it turned to be small in the zero temperature approximation. At finite temperature the (p-p), (h-h) couplings or the couplings arisen by the poles at energies $\pm(\omega_{\lambda_1 i_1} - \omega_{\lambda_2 i_2})$ can lead to an influence on the strength distribution for electromagnetic transitions.

3. CONCLUSION

We have developed the QPNM to describe correctly the interaction with (2p-2h) configurations at finite temperature by taking exactly into account the occupation numbers of one-phonon levels and obtained explicitly the relevant system of QPNM equations

with the effects of the phonon-ground state correlations and the phonon scattering. From this system one can infer all systems of QPNM equations at zero temperature obtained in Refs. /14,16,17,27/ as special cases. As has been mentioned, the tranformation from the QPNM diagrams into the NFT and TFFS diagrams takes place, which means these models could be considered equivalent. It was shown that additional poles appeared at energies $\pm(\omega_{\lambda_1 i_1} - \omega_{\lambda_2 i_2})$, which led to noncollective excitations corresponding to the phonon scattering. Our evaluation shows, that at low temperatures (T < 1 MeV) the contribution of these poles can be neglected and in this temperature region the zero temperature limit of the QPNM is a good approximation. The system of equations we have obtained can be employed to calculate the fragmentation of collective modes for even-even spherical nuclei.

The author would like to thank Professor V.G.Soloviev and Dr.A.I.Vdovin for useful discussions and support. He is grateful to Dr.G.Kyrchev for help and criticism.

REFERENCES

1. Newton J.O. et al. Phys.Rev.Lett., 1981, 46, No.21, p.1383.
2. Draper J.E. et al. Phys.Rev.Lett., 1982, 49, No.7, p.434.
3. Gobbi A., Nörenberg W. Heavy Ion Collisions, 1980, vol.2, ed. Bock R., North Holland, Amsterdam.
4. Ignatyuk A.V. Izv.Acad.Nauk SSSR, ser.fiz., 1974, 38, No.12, p.2613.
5. Ignatyuk A.V. Yad.Fiz., 1975, 21, No.1, p.21.
6. Baznat M.I. Ignatyuk A.V., Pyatov N.I. Yad.Fiz., 1979, 30, No.4(10), p. 949.
7. Ignatyuk A.V. Statistical Properties of Excited Atomic Nuclei. Energoatomizdat, M., 1983.
8. Michael Sommermann H., Ann.Phys., 1983, 151, p.163.
9. Vautherin D., Vinh Mau N. NPNO, preprint F.91406, Orsay,1983.
10. Ngo H. Thèse de Doctorat d'Etat, No.2776, Orsay, 1983.
11. Meyer F., Quentin P., Brack M. Phys.Lett., 1983, B133, No.5, p.279.
12. Brack M., Quentin P. Phys.Lett., 1974, 52B, No.2, p.159.
13. Levit S., Alhassid Y. Nucl.Phys., 1984, A413, No.3, p.439.
14. Soloviev V.G. Theory of Complex Nuclei, Pergamon Press, Oxford, 1976; Soloviev V.G. Nucleonika, 1978, vol.23,p.1149; Soloviev V.G. Sov.Phys.Particles and Nuclei, 1978, 9, No.4, p.580.
15. Nguyen Dinh Dang et al. JINR, E4-83-680, Dubna, 1983.
16. Nguyen Dinh Dang, Ponomarev V.Yu. Teor.Mat.Fiz., 1983, 57, No.1, p.154.

17. Nguyen Dinh Dang, Voronov V.V. Izv.Acad.Nauk SSSR, ser.fiz., 1984, 48, No.5, p.857.
18. Bertsch G.F., Bortignon P.F., Broglia R.A. Rev.Mod.Phys., 1983, 55, No.1, p.287.
19. Bortignon P.F., Broglia R.A. Nucl.Phys., 1981, A371, No.3, p.405.
20. Bertsch G.F. et al. Phys.Lett., 1979, 80B, No.3, p.161.
21. Bes D.R. Suppl.Prog.Theor.Phys., 1983, 74/75, p.1.
22. Vambach J., Mishra V.K., Li Chu-Hsia, Nucl.Phys., 1982, A380, No.2, p.285.
23. Kamerdziev S.P. Sov.Phys.JETPh Lett., 1979, 30, No.8,p.532; Kamerdziev S.P., Tkachev V.N. Yad.Fiz., 1982, 36, No.1(7), p.72; Kamerdziev S.P., Tseliaev V.I. Izv.Acad.Nauk SSSR, ser.fiz., 1983, 47, No.5, p.917; Kamerdziev S.P. Yad.Fiz., 1983, 38, No.2(8), p.316.
24. Stoyanov Ch., Yudin I.P. JINR, P4-11076, Dubna, 1977.
25. Beliaev S.T., Zelevinsky V.G. Nucl.Phys., 1962, 39, No.4, p.582; Marumori T., Yamamura M., Tokunaga A. Prog.Theor.Phys., 1964, 31, p.1109; Yamamura M. Prog.Theor.Phys., 1965, 33, p.199; Janssen D. et al. Nucl.Phys., 1971, A172, No.1,p.145; Marshalek E.R. Suppl.Prog.Theor.Phys., 1983, 74/75, p.89; Yamamura M. Suppl.Prog.Theor.Phys., 1983, 74/75, p.271; Tamura T., Kishimoto T. Suppl.Prog.Theor.Phys., 1983, 74/75, p.282.
26. Ponomarev V.Yu., Stoyanova O., Stoyanov Ch. JINR, P4-81-704, Dubna, 1981.
27. Kyrchev G. Teor.Mat.Fiz., 1982, 53, No.2, p.260.
28. Zubarev D.N. Usp.Fiz.Nauk, 1960, 71, No.1, p.71.
29. Abricosov A.A., Gorkov A.P., Dzialoshinsky I.E. Method of Quantum Field Theory in Statistical Physics, Fizmatgiz, M., 1962.
30. Yamasaki S. Prog.Theor.Phys., 1983, 70, No.4, p.995.
31. Lipkin H.J. Quantum Mechanics. New Appoaches to Selected Topics. Amsterdam-London, North-Holland, New York, Elsevier, 1973.
32. Soloviev V.G., Stoyanova O., Stoyanov Ch. Izv.Acad.Nauk SSSR, ser.fiz., 1980, 44, No.9, p.1938; Navrotska-Rybarska V., Stoyanova O., Stoyanova Ch. Yad.Fiz., 1981, 33, No.6, p.1494; Stoyanova O. JINR, P4-81-477, Dubna, 1981.
33. Nguyen Dinh Dang, Soloviev V.G. JINR, P4-83-325, Dubna, 1983.

Received by Publishing Department
on July 6, 1984.

Нгуен Динь Данг E4-84-481

Система уравнений квазичастично-фононной модели ядра с учетом рассеяния фононов при конечной температуре

Дано обобщение квазичастично-фононной модели ядра /КФМ/ для описания взаимодействия с (2p-2h) конфигурациями при конечной температуре. Путем точного учета чисел заполнения однофононных уровней получена в явном виде замкнутая система уравнений для функции Грина с пропагаторами, соответствующими переходам из одного в два фонона, фононным корреляциям в основном состоянии и фононному рассеянию в четно-четных сферических ядрах. Установлено взаимно-однозначное соответствие между этой системой и системой уравнений КФМ для коэффициентов волновой функции возбужденного состояния. Показано, что в пределе нулевой температуры можно получить основные уравнения, использованные до сих пор в КФМ. Численная оценка эффектов фононного рассеяния показала, что при температурах $T < 1$ МэВ хорошо применим предел нулевой температуры КФМ. Проведено сравнение диаграмм КФМ с диаграммами теории ядерных полей и результатами теории конечных Ферми-систем.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна 1984

Nguyen Dinh Dang E4-84-481

Equations of the Quasiparticle-Phonon Nuclear Model with the Phonon Scattering Effects at Finite Temperature

A generalization of the quasiparticle-phonon nuclear model /QPNM/ for describing the interaction with (2p-2h) configurations at finite temperature is presented. By taking exactly into account the occupation numbers of one-phonon energy levels a closed system of approximate equations for Green functions with one- to two-phonon transition, phonon-ground state correlation and phonon scattering propagators in even-even spherical nuclei is explicity derived. A one-to-one correspondence between this system and the system of QPNM equations of the coefficients of the excited state wave function is established. It is shown that in the zero temperature limit one obtains the standard basic equations employed so far within the QPNM. The numerical evaluation of the phonon scattering effects has shown that for temperatures $T < 1$ MeV the zero temperature limit of the QPNM is a quite good approximation. The equivalence between the QPNM diagrams and the diagrams of the nuclear field theory and the theory of finite Fermi-systems is discussed.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

Preprint of the Joint Institute for Nuclear Research. Dubna 1984