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# BASIC EQUATIONS FOR ODD SPHERICAL NUCLEI IN THE QUASIPARTICLE-PHONON NUCLEAR MODEL

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In recent years the quasiparticle-phonon nuclear model (QPM)[1-4] is widely used for the description of fragmentation of nuclear excitations of simple structure (such as one-phonon or one-quasiparticle states) at intermediate and high excitation energies. Within this model one is able to describe a wide set of experimental data in spectroscopy for even-even[2] and odd[3,4] spherical nuclei.

Despite the successes of the QPM, a further improvement of the model is necessary due to the fact that some physical effects firstly not included into the model formalism could play an important role in the number of cases. Such investigations have been performed in[5-7] for even-even nuclei. In this paper we consider odd spherical nuclei.

The QPM equations for odd spherical nuclei have been obtained in[8], where the considered model wave functions contain components ranging in complexity from one quasiparticle to "quasiparticle plus three phonons". In the equations used in numerical calculations only contributions from components not more complex than "quasiparticle plus two phonons" are taken into account[3]. However, in the equations from refs.[8] and [3], the anharmonic effects of phonon excitations of even-even core and Pauli principle corrections due to the fermion structure of the phonon state are not taken into account.

The influence of the mentioned effects on the fragmentation of highly excited single-particle modes, which has been analysed by the use of equations from[3,8], is expected to be not strong. But, as has been shown by investigations in[4,9], these effects are substantial at low excitation energies and influence strongly the fragmentation of "quasiparticle plus two phonons" states.

In this paper, within the QPM, we obtain the equations for the energies of excited states and structure coefficients in their wave functions using model wave function of an odd spherical nucleus in the form:

$$\Psi_\nu(JM) = C_{J\nu}\{\alpha^+_{JM} + \sum_{\lambda i j} D_j^{\lambda i}(J\nu)[\alpha^+_{jm} Q^+_{\lambda\mu i}]_{JM} + \sum_{\substack{\lambda_1 i_1 \lambda_2 i_2 \\ \lambda j}} F_{j\lambda}^{\lambda_1 i_1 \lambda_2 i_2}(J\nu)[\alpha^+_{jm}[Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2}]_{\lambda\mu}]_{JM}\}\Psi_0 , \tag{1}$$



where $\alpha^+_{jm}$ and $Q^+_{\lambda\mu i}$ are the creation operators of Bogolubov's quasiparticle with quantum numbers $j \equiv (n, \ell, j)$ and projection of total moment $m$, and $i$-th phonon with moment and projection $\lambda\mu$, respectively. $\Psi_0$ is the ground state wave function of even-even core. The phonon creation operator is expressed in terms of quasiparticle creation and annihilation operators as follows:

$$Q^+_{\lambda\mu i} = \frac{1}{2}\sum_{jj'} \psi^{\lambda i}_{jj'}[\alpha^+_{jm}\alpha^+_{j'm'}]_{\lambda\mu} + (-)^{\lambda-\mu}\phi^{\lambda i}_{jj'}[\alpha_{j'm'}\alpha_{jm}]_{\lambda-\mu} . \tag{2}$$

We write the commutation relations between phonon and quasiparticle operators:

$$[Q_{\lambda_1\mu_1 i_1}, Q^+_{\lambda_2\mu_2 i_2}] = \delta_{\lambda_1\lambda_2}\delta_{\mu_1\mu_2}\sum_{jj'}\frac{1}{2}(\psi^{\lambda_1 i_1}_{jj'}\psi^{\lambda_2 i_2}_{jj'} - \phi^{\lambda_1 i_1}_{jj'}\phi^{\lambda_2 i_2}_{jj'}) - \sum_{jj'j_1}\sum_{mm'm_1}\{\psi^{\lambda_1 i_1}_{j'j_1}\psi^{\lambda_2 i_2}_{jj_1}\langle j'm'j_1m_1|\lambda_1\mu_1\rangle\langle jmj_1m_1|\lambda_2\mu_2\rangle - (-)^{\lambda_1+\lambda_2-\mu_1-\mu_2}\phi^{\lambda_1 i_1}_{jj_1}\phi^{\lambda_2 i_2}_{j'j_1}\langle jmj_1m_1|\lambda_1-\mu_1\rangle\langle j'm'j_1m_1|\lambda_2-\mu_2\rangle\}\alpha^+_{jm}\alpha_{j'm'} , \tag{3}$$

$$[\alpha^+_{j_1m_1}, Q_{\lambda\mu i}] = [Q^+_{\lambda\mu i}, \alpha_{j_1m_1}]^+ = \sum_{jm}\langle jmj_1m_1|\lambda\mu\rangle\psi^{\lambda i}_{jj_1}\alpha_{jm} . \tag{4}$$

For nuclei with negligible number of quasiparticles in the ground state, i.e., when the random phase approximation (RPA) is valid, one has $\langle\Psi_0\,\alpha^+_{jm}\alpha_{jm}\Psi_0\rangle \sim 0$, so from the orthonormalization relation of one-phonon states one gets

$$\sum_{jj'}\psi^{\lambda_1 i_1}_{jj'}\psi^{\lambda_2 i_2}_{jj'} - \phi^{\lambda_1 i_1}_{jj'}\phi^{\lambda_2 i_2}_{jj'} = 2\delta_{\lambda_1\lambda_2}\delta_{i_1 i_2} . \tag{5}$$

Further, we suppose that the $\psi$ and $\phi$ amplitudes satisfy the RPA equations for the separable interaction between quasiparticles and thus the condition (5).

The QPM Hamiltonian consists of phenomenological mean fields for neutrons and protons, usual monopole pairing interaction in the particle-particle channel and separable multipole and spin-multipole forces in the particle-hole channel. Interaction in the particle-hole channel is isotopically invariant. A more detailed description of Hamiltonian can be found, for example in[2,10].

We write the QPM Hamiltonian in terms of quasiparticle and phonon operators, assuming that $\psi$ and $\phi$ amplitudes do satisfy the RPA equations.

$$H = \sum_{jm} \epsilon_j \alpha^+_{jm}\alpha_{jm} - \frac{1}{4}\sum_{\lambda\mu i i' \tau} \frac{X^{\lambda i}(\tau) + X^{\lambda i'}(\tau)}{\sqrt{\mathcal{Y}^{\lambda i}_\tau \mathcal{Y}^{\lambda i'}_\tau}} \times$$

$$\times (Q^+_{\lambda\mu i} + (-)^{\lambda-\mu} Q_{\lambda-\mu i})(Q^+_{\lambda-\mu i'}(-)^{\lambda-\mu} + Q_{\lambda\mu i'}) - \qquad (6)$$

$$- \frac{1}{2\sqrt{2}}\sum_{\lambda\mu i}\{(Q^+_{\lambda\mu i}(-)^{\lambda-\mu} + Q_{\lambda-\mu i})\sum_{jj'\tau}\frac{f^{\lambda}_{jj'} v^{(-)}_{jj'}}{\sqrt{\mathcal{Y}^{\lambda i}_\tau}} B(jj';\lambda-\mu) + h.c.\}$$

we have denoted

$$X^{\lambda i}(\tau) = \frac{1}{2\lambda+1}\sum^{\tau}_{jj'} \frac{(f^{\lambda}_{jj'} u^{(+)}_{jj'})^2 \epsilon_{jj'}}{\epsilon^2_{jj'} - \omega^2_{\lambda i}},$$

$$\mathcal{Y}^{\lambda i}_\tau = Y^{\lambda i}_\tau + Y^{\lambda i}_{-\tau}\left\{\frac{1-(\kappa^{(\lambda)}_0 + \kappa^{(\lambda)}_1) X^{\lambda i}(\tau)}{(\kappa^{(\lambda)}_0 - \kappa^{(\lambda)}_1) X^{\lambda i}(-\tau)}\right\}^2,$$

$$Y^{\lambda i}_\tau = \frac{1}{2}\frac{\partial}{\partial\omega} X^{\lambda i}(\tau)|_{\omega=\omega_{\lambda i}}; \quad \epsilon_{jj'} = \epsilon_j + \epsilon_{j'},$$

where $\epsilon_j$ is quasiparticle energy, $f^{\lambda}_{jj'} = \langle j|| R_\lambda(r) i^\lambda Y_{\lambda\mu} ||j'\rangle$ $u^{(+)}_{jj'} = u_j v_{j'} + u_{j'} v_j$, $v^{(-)}_{jj'} = u_j u_{j'} - v_j v_{j'}$ ($u_j$, $v_j$ are Bogolubov's transformation coefficients), $\tau$ is isotopic index: $\tau = (n, p)$. Sign change $\tau \to -\tau$ corresponds to the interchange $n \leftrightarrow p$. Index $\tau$ in $\Sigma^\tau$ means that the sum runs over only one type of single-particle spectrum: neutron or proton. Using (3) and (4), we get the normalization condition of the wave function (4) in the form:

$$\langle \Psi^*_\nu(JM)\Psi_\nu(JM)\rangle = 1 = C^2_{J\nu}\{1 + \sum_{\lambda ij}[D^{\lambda i}_j(J\nu)]^2 +$$

$$+ \sum_{12}\mathcal{L}_J(j_1\lambda_1 i_1|j_2\lambda_2 i_2) D^{\lambda_1 i_1}_{j_1}(J\nu) D^{\lambda_2 i_2}_{j_2}(J\nu) + 2\sum_{123}[F^{\lambda_1 i_1\lambda_2 i_2}_{j_3\lambda_3}(J\nu)]^2 +$$

$$+ \sum_{1231'2'3'} F^{\lambda_1 i_1\lambda_2 i_2}_{j_3\lambda_3}(J\nu) F^{\lambda'_1 i'_1\lambda'_2 i'_2}_{j_3\lambda'_3}(J\nu) \mathcal{K}^J_{j_3}(\lambda'_3, \lambda'_2 i'_2, \lambda'_1 i'_1|\lambda_1 i_1, \lambda_2 i_2, \lambda_3) +$$

$$+ \sum_{1231'3'} F^{\lambda_1 i_1\lambda_2 i_2}_{j_3\lambda_3}(J\nu) F^{\lambda'_1 i'_1\lambda_2 i_2}_{j'_3\lambda'_3}(J\nu) \mathfrak{M}^J_{\lambda_2}(j'_3\lambda'_3, \lambda'_1 i'_1|\lambda_1 i_1, j_3\lambda_3) +$$



$$+ \sum_{1232'3'} F^{\lambda_2 i_2\lambda_1 i_1}_{j_3\lambda_3}(J\nu) F^{\lambda'_2 i'_2\lambda_1 i_1}_{j'_3\lambda'_3}(J\nu) \mathfrak{M}^J_{\lambda_1}(j'_3\lambda'_3, \lambda'_2 i'_2|\lambda_2 i_2, j_3\lambda_3) +$$

$$+ \frac{1}{2}\sum_{\substack{1234\\1'2'3'4'}} F^{\lambda_1 i_1\lambda_2 i_2}_{j_3\lambda_3}(J\nu) F^{\lambda'_1 i'_1\lambda'_2 i'_2}_{j'_3\lambda'_3}(J\nu) \times \qquad (7)$$

$$\times [\mathcal{K}^J_{j_3}(\lambda'_4, \lambda'_2 i'_2, \lambda_4 i_4|\lambda_1 i_1, \lambda_2 i_2, \lambda_3) \mathfrak{M}^J_{\lambda'_2}(j'_3\lambda'_3, \lambda'_1 i'_1|\lambda_4 i_4, j_3\lambda'_4) +$$

$$+ \mathcal{K}^J_{j_3}(\lambda'_4, \lambda_4 i_4, \lambda'_1 i'_1|\lambda_1 i_1, \lambda_2 i_2, \lambda_3) \mathfrak{M}^J_{\lambda'_1}(j'_3\lambda'_3, \lambda'_2 i'_2|\lambda_4 i_4, j_3\lambda'_4)]\},$$

where $\mathcal{L}_J$, $\mathfrak{M}^J_\lambda$ and $\mathcal{K}^J_j$ functions have forms:

$$\mathcal{L}_J(j_1\lambda_1 i_1|j_2\lambda_2 i_2) = \sum_{j_3}\hat{\lambda}_1\hat{\lambda}_2\begin{Bmatrix} j_2 & j_3 & \lambda_1\\ j_1 & J & \lambda_2\end{Bmatrix}\psi^{\lambda_1 i_1}_{j_3 j_2}\psi^{\lambda_2 i_2}_{j_3 j_1},$$

$$\mathcal{K}^J_j(\lambda', \lambda'_2 i'_2, \lambda'_1 i'_1|\lambda_1 i_1, \lambda_2 i_2, \lambda) = \sum_{\substack{\mu_1\mu_2\mu'_1\mu'_2\\ m\mu\mu'}}\langle jm\lambda\mu|JM\rangle\langle jm\lambda'\mu'|JM\rangle \times$$

$$\times \langle\lambda_1\mu_1\lambda_2\mu_2|\lambda\mu\rangle\langle\lambda'_1\mu'_1\lambda'_2\mu'_2|\lambda'\mu'\rangle\, \mathcal{K}(\lambda'_2\mu'_2 i'_2, \lambda'_1\mu'_1 i'_1|\lambda_1\mu_1 i_1, \lambda_2\mu_2 i_2),$$

$$\mathfrak{M}^J_{\lambda_2}(j'\lambda', \lambda'_1 i'_1|\lambda_1 i_1, j\lambda) = -2\hat{\lambda}\hat{\lambda}'\hat{\lambda}_1\hat{\lambda}'_1\sum_{j_1}\begin{Bmatrix} J & \lambda' & j'\\ \lambda & \lambda_2 & \lambda_1\\ j & \lambda'_1 & j_1\end{Bmatrix}\psi^{\lambda'_1 i'_1}_{j j_1}\psi^{\lambda_1 i_1}_{j' j_1}. \qquad (8)$$

With $\hat{\lambda} = \sqrt{2\lambda+1}$; function $\mathcal{K}(\lambda'_2\mu'_2 i'_2, \lambda'_1\mu'_1 i'_1|\lambda_1\mu_1 i_1, \lambda_2\mu_2 i_2)$ is a combination of $\psi$ and $\phi$-phonon amplitudes, and its explicit form is given in [6].

The average of Hamiltonian H over the wave functions (1) is equal to:

$$\langle\Psi^*_\nu(JM) H\Psi_\nu(JM)\rangle = C^2_{J\nu}\{\epsilon_J + \sum_{\lambda ij}[D^{\lambda i}_j(J\nu)]^2(\epsilon_j + \omega_{\lambda i}) +$$

$$+ \frac{1}{2}\sum_{12} D^{\lambda_1 i_1}_{j_1}(J\nu) D^{\lambda_2 i_2}_{j_2}(J\nu)[\mathcal{L}_J(j_1\lambda_1 i_1|j_2\lambda_2 i_2)(\epsilon_{j_1 j_2} + \omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2}) -$$

$$- R_J(j_1\lambda_1 i_1|j_2\lambda_2 i_2)] - \sqrt{2}\sum_{\lambda ij} D^{\lambda i}_j(J\nu)[\Gamma(Jj\lambda i) +$$

$$+ \sum_{\lambda' i' j'}\Gamma(Jj'\lambda' i')\mathcal{L}_J(j\lambda i|j'\lambda' i')] + 2\sum_{123}\{[F^{\lambda_1 i_1\lambda_2 i_2}_{j_3\lambda_3}(J\nu)]^2 \times$$

$$\times (\epsilon_j + \omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2}) + \sum_{1'2'3'} F^{\lambda_1 i_1\lambda_2 i_2}_{j_3\lambda_3}(J\nu) F^{\lambda'_1 i'_1\lambda'_2 i'_2}_{j'_3\lambda'_3}(J\nu) \times$$

$$\times[\frac{1}{2}K^J_{j_3}(\lambda'_3,\lambda'_2 i'_2,\lambda'_1 i'_1|\lambda_1 i_1,\lambda_2 i_2,\lambda_3)\delta_{j_3 j'_3} + \mathfrak{M}^J_{\lambda_2}(j'_3\lambda'_3,\lambda'_1 i'_1|\lambda_1 i_1, j_3\lambda_3)\delta_{\lambda_2\lambda'_2}\delta_{i_2 i'_2} +$$

$$+\frac{1}{2}\sum_{45} K^J_{j_3}(\lambda_4,\lambda'_2 i'_2,\lambda_5 i_5|\lambda_1 i_1,\lambda_2 i_2,\lambda_3)\mathfrak{M}^J_{\lambda'_2}(j'_3\lambda'_3,\lambda'_1 i'_1|\lambda_5 i_5, j_3\lambda_4)]\times$$

$$\times(\epsilon_{j_3}+\omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2}) - S_J(j'_3\lambda'_3,\lambda'_2 i'_2,\lambda'_1 i'_1|\lambda_1 i_1,\lambda_2 i_2, j_3\lambda_3)\} +$$

$$+2\sum_{\substack{123\\1'2'3'}} F^{\lambda_1 i_1\lambda_2 i_2}_{j_3\lambda_3}(J\nu) D^{\lambda_3 i_3}_{j_3}(J\nu)[\delta_{j_3 j'_3}\delta_{\lambda_3\lambda'_3}\delta_{i_3 i'_3} + \mathfrak{L}_J(j'_3\lambda'_3 i'_3|j_3\lambda_3 i_3)]\times$$

$$\times[U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda'_3 i'_3) + \frac{1}{2}U^{\lambda'_2 i'_2}_{\lambda'_1 i'_1}(\lambda'_3 i'_3) K^J_{j'_3}(\lambda'_3,\lambda'_2 i'_2,\lambda'_1 i'_1|\lambda_1 i_1,\lambda_2 i_2,\lambda_3)] +$$

$$+2\sqrt{2}\sum_{\substack{1234\\1'2'3'4'}} F^{\lambda_1 i_1\lambda_2 i_2}_{j'_3\lambda'_3}(J\nu) D^{\lambda_3 i_3}_{j_3}(J\nu)\hat{j}_4\hat{\lambda}_4(-)^{j'_3+\lambda'_2+\lambda'_1-J}\begin{Bmatrix}\lambda'_1 & \lambda'_2 & \lambda_4\\ J & j'_3 & j_4\end{Bmatrix}\times$$

$$\times[\delta_{j_4 j_3}\delta_{\lambda_2\lambda_3}\delta_{i_2 i_3}\delta_{\lambda'_1\lambda_1}\delta_{i'_1 i_1}\delta_{i'_2 i_2}\delta_{\lambda'_2\lambda_2}\delta_{\lambda_4\lambda'_3}(\Gamma(j_4 j'_3\lambda'_1 i'_1) +$$

$$+\sum_{4'}\Gamma(j_3 j'_4\lambda'_4 i'_4)\mathfrak{L}_J(j'_4\lambda'_4 i'_4|j'_3\lambda_1 i_1)) + \Gamma(j_4 j'_3\lambda'_1 i'_1)\times$$

$$\times(\delta_{\lambda'_1\lambda_1}\delta_{i'_1 i_1}\delta_{\lambda'_2\lambda_2}\delta_{i'_2 i_2}\delta_{\lambda_4\lambda'_3}\mathfrak{L}_J(j_4\lambda_2 i_2|j_3\lambda_3 i_3) +$$

$$+\frac{1}{2}K^J_{j'_3}(\lambda_4,\lambda_3 i_3,\lambda'_1 i'_1|\lambda_1 i_1,\lambda_2 i_2,\lambda'_3)(\delta_{j_3 j_4}\delta_{\lambda'_2\lambda_3}\delta_{i'_2 i_3} +$$

$$+\mathfrak{L}_J(j_4\lambda'_2 i'_2|j_3\lambda_3 i_3))) + \frac{1}{2}\sum_{4'}\Gamma(j_3 j'_4\lambda'_4 i'_4)\times \tag{9}$$

$$\times K^J_{j'_3}(\lambda_4,\lambda_3 i_3,\lambda'_4 i'_4|\lambda_1 i_1,\lambda_2 i_2,\lambda'_3)\mathfrak{L}_J(j'_4\lambda'_1 i'_1|j'_3\lambda_1 i_1)\delta_{\lambda'_2\lambda_2}\delta_{i'_2 i_2}\delta_{j_3 j_4}]\},$$

where $\Gamma(Jj\lambda i) = \frac{\hat{\lambda}}{\hat{f}}\frac{f^{\lambda}_{Jj}\, v^{(-)}_{Jj}}{\sqrt{\mathcal{Y}^{\lambda i}_r}}$,

$$R_J(j_1\lambda_1 i_1|j_2\lambda_2 i_2) = \sum_{i'}^{r}\{X^r(\lambda_2 i_2,\lambda_2 i')\mathfrak{L}_J(j_1\lambda_1 i_1|j_2\lambda_2 i') +$$

$$+X^r(\lambda_1 i_1,\lambda_1 i')\mathfrak{L}_J(j_2\lambda_2 i_2|j_1\lambda_1 i') + \sum_{j_3\lambda i} X^r(\lambda i,\lambda i')\times$$

$$\times[\mathfrak{L}_J(j_1\lambda_1 i_1|j_3\lambda i)\mathfrak{L}_J(j_2\lambda_2 i_2|j_3\lambda i') +$$

$$+\mathfrak{L}_J(j_2\lambda_2 i_2|j_3\lambda i)\mathfrak{L}_J(j_1\lambda_1 i_1|j_3\lambda i')]\},$$



$$X^r(\lambda_1 i_1,\lambda_2 i_2) = \frac{1}{4}\frac{X^{\lambda_1 i_1}(r) + X^{\lambda_2 i_2}(r)}{\sqrt{\mathcal{Y}_r^{\lambda_1 i_1}\mathcal{Y}_r^{\lambda_2 i_2}}}, \tag{10}$$

$$S_J(j'_3\lambda'_3,\lambda'_2 i'_2,\lambda'_1 i'_1|\lambda_1 i_1,\lambda_2 i_2,j_3\lambda_3) = -\frac{1}{2}\sum_{44'}^{r}\{[\mathfrak{M}^J_{\lambda_2}(j'_3\lambda'_3,\lambda'_1 i'_1|\lambda_4 i_4, j_3\lambda'_4) +$$

$$+\delta_{j'_3 j_3}\delta_{\lambda_4\lambda'_1}\delta_{i_4 i'_1}\delta_{\lambda'_3\lambda'_4}][X^r(\lambda_4 i_4,\lambda_4 i'_4)K^J_{j_3}(\lambda'_4,\lambda'_2 i'_2,\lambda_4 i'_4|\lambda_1 i_1,\lambda_2 i_2,\lambda_3) +$$

$$+X^r(\lambda'_2 i'_2,\lambda'_2 i'_4)K^J_{j_3}(\lambda'_4,\lambda'_2 i'_4,\lambda_4 i_4|\lambda_1 i_1,\lambda_2 i_2,\lambda_3) + \sum_{567}^{r} X^r(\lambda_5 i_5,\lambda_5 i'_4)\times$$

$$\times K^J_{j_3}(\lambda_7,\lambda_6 i_6,\lambda_5 i'_4|\lambda_1 i_1,\lambda_2 i_2,\lambda_3)K^J_{j_3}(\lambda'_4,\lambda'_2 i'_2,\lambda_4 i_4|\lambda_5 i_5,\lambda_6 i_6,\lambda_7)] -$$

$$-\sum_{56}^{r} X^r(\lambda_4 i_4,\lambda_4 i'_4)\mathfrak{M}^J_{\lambda_2}(j_3\lambda_3,\lambda_1 i_1|\lambda_4 i'_4, j_3\lambda'_4)\times$$

$$\times[\mathfrak{M}^J_{\lambda_2}(j_5\lambda'_4,\lambda_4 i_4|\lambda_5 i_5, j'_3\lambda_6)(\delta_{\lambda_5\lambda'_1}\delta_{i_5 i'_1}\delta_{\lambda_2\lambda'_2}\delta_{i_2 i'_2}\delta_{\lambda'_3\lambda_6} +$$

$$+\frac{1}{2}K^J_{j'_3}(\lambda'_3,\lambda'_2 i'_2,\lambda'_1 i'_1|\lambda_5 i_5,\lambda_2 i_2,\lambda_6)) + (-)^{\lambda'_4+\lambda_6+\lambda_2+\lambda_5}\times$$

$$\times\mathfrak{M}^J_{\lambda_4}(j_5\lambda'_4,\lambda_2 i_2|\lambda_5 i_5, j'_3\lambda_6)(\delta_{\lambda_4\lambda'_1}\delta_{i_4 i'_1}\delta_{\lambda_5\lambda'_2}\delta_{i_5 i'_2}\delta_{\lambda'_3\lambda_6} +$$

$$+\frac{1}{2}K^J_{j'_3}(\lambda'_3,\lambda'_2 i'_2,\lambda'_1 i'_1|\lambda_4 i_4,\lambda_5 i_5,\lambda_6))] - 2(\delta_{\lambda_4\lambda'_1}\delta_{i_4 i'_1}\delta_{\lambda_2\lambda'_2}\delta_{i_2 i'_2}\delta_{\lambda'_3\lambda'_4} +$$

$$+\frac{1}{2}K^J_{j'_3}(\lambda'_3,\lambda'_2 i'_2,\lambda'_1 i'_1|\lambda_4 i_4,\lambda_2 i_2,\lambda'_4))X^r(\lambda_4 i_4,\lambda_4 i'_4)\times$$

$$\times\mathfrak{M}^J_{\lambda_2}(j_3\lambda_3,\lambda_1 i_1|\lambda_4 i'_4, j'_3\lambda'_4)\}. \tag{11}$$

The expression for matrix elements $U^{\lambda_1 i_1}_{\lambda_2 i_2}(\lambda i)$ of Hamiltonian (6) between one- and two-phonon states is given in [2]. The inclusion of terms containing $U^{\lambda_1 i_1}_{\lambda_2 i_2}(\lambda i)$ into (9) accounts for the anharmonics of vibrations of even-even core. Using variational principle we could obtain the general system of QPM equations with the wave function (1). But, before going to do this, we make some approximations for (7) and (9). $\mathfrak{L}_J$, $K^J_j$ and $\mathfrak{M}^J_\lambda$ functions appear with taking into account Pauli principle in the wave function (1). If one neglects Pauli principle corrections these functions are equal to zero, As has been shown in [6,7,9], $\mathcal{L}_J$ and $K^J_j$ functions are alternating quantities and their diagonal values are much greater than nondiagonal ones. Our investigations show a similar situation for $\mathfrak{M}^J_\lambda$ functions. It can easily be seen from the table, where $\mathfrak{M}^J_\lambda$ values with $\lambda = 2$ for states with $J^\pi = 9/2^+$ in $^{209}$Pb and $J^\pi = 1/2^-$ in $^{61}$Ni are shown. Values of $\mathfrak{M}^J_\lambda$, as in the case for $\mathcal{L}_J$ and $K^J_j$, have different signs and reach maximum in diagonal terms.

The maximum of $\mathcal{L}_J$, $\mathfrak{M}^J_\lambda$ and $K^J_j$ functions in their diagonal values is natural from the physical point of view, because Pauli principle is violated most probably in the configurations formed by identical quasiparticles. Keeping this in mind, we only use diagonal approximation in our further calculations. In this case (7) becomes

$$<\Psi^*_\nu(JM)\Psi_\nu(JM)> = 1 = C^2_{J\nu}\{1 + \sum_{\lambda ij}[D^{\lambda i}_j(J\nu)]^2[1 + \mathcal{L}(Jj\lambda i)] +$$

$$+ 2\sum_{123}[F^{\lambda_1 i_1 \lambda_2 i_2}_{j_3\lambda_3}(J\nu)]^2[1 + \frac{1}{2}K^{\lambda_3}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)] \times \tag{12}$$

$$\times [1 + \mathfrak{M}(Jj_3\lambda_3|\lambda_1 i_1, \lambda_2 i_2)]\},$$

where $\mathcal{L}(Jj\lambda i) = \mathcal{L}_J(j\lambda i|j\lambda i)$,

$$\mathfrak{M}(Jj\lambda|\lambda_1 i_1, \lambda_2 i_2) = \frac{1}{2}[\mathfrak{M}^J_{\lambda_2}(j\lambda, \lambda_1 i_1|\lambda_1 i_1, j\lambda) + \mathfrak{M}^J_{\lambda_1}(j\lambda, \lambda_2 i_2|\lambda_2 i_2, j\lambda)],$$

$$K^\lambda(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2) = \sum_{1234}(-)^{j_2+j_4-\lambda}(2\lambda_1+1)(2\lambda_2+1)\begin{Bmatrix} j_1 & j_2 & \lambda_2 \\ j_4 & j_3 & \lambda_1 \\ \lambda_1 & \lambda_2 & \lambda \end{Bmatrix} \times$$

$$\times [\psi^{\lambda_1 i_1}_{j_3 j_4}\psi^{\lambda_1 i_1}_{j_1 j_4}\psi^{\lambda_2 i_2}_{j_3 j_2}\psi^{\lambda_2 i_2}_{j_1 j_2} - \phi^{\lambda_1 i_1}_{j_3 j_4}\phi^{\lambda_1 i_1}_{j_1 j_4}\phi^{\lambda_2 i_2}_{j_3 j_2}\phi^{\lambda_2 i_2}_{j_1 j_2}].$$

For the "quasiparticle plus phonon" configurations, which are forbidden by Pauli principle, $\mathcal{L}(Jj\lambda i) = -1$ and these states automatically are excluded from the normalization condition and following expressions [9]. If violation of Pauli principle takes place in the two phonon components, $K^\lambda = -2$ and these components are also excluded from consideration [6,7]. At last, if Pauli principle is violated in the "quasiparticle plus two phonon" component, namely, between quasiparticle and one phonon, $\mathfrak{M} = -1$ and this component is excluded automatically. For example, if two-quasiparticle component $\{2p_{3/2} \otimes 2p_{3/2}\}_{2^+}$ gives 100% contribution to the structure of $2^+_2$ phonon in $^{60}$Ni for $1/2^-$ state in $^{61}$Ni which is forbidden by Pauli principle,



$$\mathfrak{M}(\tfrac{1}{2}\,\tfrac{3}{2}2|4^+, 2^+_2) = \frac{1}{2}\mathfrak{M}^{1/2}_4(\tfrac{3}{2}2, 2^+_2|2^+_2, \tfrac{3}{2}2) = -1. \tag{13}$$

Thus, the strict commutation relations between phonon and quasiparticle operators enable one to take into account Pauli principle.

In the diagonal approximation the average of Hamiltonian is much simplified and has the following form:

$$<\Psi^*_\nu(JM)H\Psi_\nu(JM)> = C^2_{J\nu}\{\epsilon_J + \sum_{\lambda ij}[(D^{\lambda i}_j(J\nu))^\lambda(\epsilon_j + \omega_{\lambda i} - R(Jj\lambda i)) -$$

$$- \sqrt{2}D^{\lambda i}_j(J\nu)\Gamma(Jj\lambda i)][1 + \mathcal{L}(Jj\lambda i)] +$$

$$+ 2\sum_{123}[F^{\lambda_1 i_1 \lambda_2 i_2}_{j_3\lambda_3}(J\nu)]^2[\epsilon_{j_3} + \omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2} - S(Jj_3\lambda_3|\lambda_1 i_1, \lambda_2 i_2)] \times$$

$$\times [1 + \frac{1}{2}K^{\lambda_3}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)][1 + \mathfrak{M}(Jj_3\lambda_3|\lambda_1 i_1, \lambda_2 i_2)] +$$

$$+ 2\sum_{123}D^{\lambda_3 i_3}_{j_3}(J\nu)F^{\lambda_1 i_1 \lambda_2 i_2}_{j_3\lambda_3}(J\nu)U^{\lambda_2 i_2}_{\lambda_1 i_1}(\lambda_3 i_3) \times$$

$$\times [1 + \frac{1}{2}K^{\lambda_3}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)][1 + \mathcal{L}(Jj_3\lambda_3 i_3)] -$$

$$- 2\sqrt{2}\sum_{123}D^{\lambda_1 i_1}_{j_1}(J\nu)F^{\lambda_2 i_2 \lambda_1 i_1}_{j_3\lambda_3}(J\nu)\hat{j}_1\hat{\lambda}_3(-)^{j_3+\lambda_2+\lambda_1+J} \times$$

$$\times \begin{Bmatrix} \lambda_2 & \lambda_1 & \lambda_3 \\ J & j_3 & j_1 \end{Bmatrix}\Gamma(j_1 j_3 \lambda_2 i_2)[1 + \frac{1}{2}K^{\lambda_3}(\lambda_1 i_1, \lambda_2 i_2|\lambda_2 i_2, \lambda_1 i_1)] \times$$

$$\times [1 + \mathcal{L}(Jj_3\lambda_2 i_2) + \mathcal{L}(Jj_1\lambda_1 i_1)]\}, \tag{14}$$

where $R(Jj\lambda i) = \dfrac{R_J(j\lambda i|j\lambda i)}{1 + \mathcal{L}(Jj\lambda i)}$,

$$S(Jj\lambda|\lambda_1 i_1, \lambda_2 i_2) = \frac{S_J(j\lambda, \lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2, j\lambda)}{[1 + \frac{1}{2}K^\lambda(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)][1 + \mathfrak{M}(Jj\lambda|\lambda_1 i_1, \lambda_2 i_2)]}.$$

From (14) one can see that states forbidden by Pauli principle give no contribution to the average of Hamiltonian. Using variational principle,

$$\delta\{<\Psi^*_\nu(JM)H\Psi_\nu(JM)> - \eta_{J\nu}[<\Psi^*_\nu(JM)\Psi_\nu(JM)> - 1]\} = 0. \tag{15}$$

we obtain the system of three equations

$$\{\epsilon_J - \eta_{J\nu} + \sum_1 [(D_{j_1}^{\lambda_1 i_1}(J\nu))^2(\epsilon_{j_1} + \omega_{\lambda_1 i_1} - R(Jj_1\lambda_1 i_1) - \eta_{J\nu}) -$$

$$-\sqrt{2}D_{j_1}^{\lambda_1 i_1}(J\nu)\Gamma(Jj_1\lambda_1 i_1)][1 + \mathcal{L}(Jj_1\lambda_1 i_1)] +$$

$$+ 2\sum_{123}[F_{j_3\lambda_3}^{\lambda_1 i_1\lambda_2 i_2}(J\nu)]^2[\epsilon_{j_3} + \omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2} - S(Jj_3\lambda_3|\lambda_1 i_1, \lambda_2 i_2) - \eta_{J\nu}] \times$$

$$\times[1 + \frac{1}{2}K^{\lambda_3}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)][1 + \mathcal{M}(Jj_3\lambda_3|\lambda_1 i_1, \lambda_2 i_2)] +$$

$$+ 2\sum_{123} D_{j_3}^{\lambda_3 i_3}(J\nu)F_{j_3\lambda_3}^{\lambda_1 i_1\lambda_2 i_2}(J\nu)U_{\lambda_1 i_1}^{\lambda_2 i_2}(\lambda_3 i_3)[1 + \frac{1}{2}K^{\lambda_3}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)] \times$$

$$\times[1 + \mathcal{L}(Jj_3\lambda_3 i_3)] - 2\sqrt{2}\sum_{123} D_{j_2}^{\lambda_2 i_2}(J\nu)F_{j_3\lambda_3}^{\lambda_1 i_1\lambda_2 i_2}(J\nu)(-)^{j_3+\lambda_1+\lambda_2+J} \times$$

$$\times \hat{j}_2\hat{\lambda}_3\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda_3\\ J & j_3 & j_2\end{Bmatrix}\Gamma(j_2 j_3\lambda_1 i_1)[1 + \frac{1}{2}K^{\lambda_3}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)] \times$$

$$\times[1 + \mathcal{L}(Jj_3\lambda_1 i_1) + \mathcal{L}(Jj_2\lambda_2 i_2)]\} = 0, \qquad (16)$$

$$\{D_j^{\lambda i}(J\nu)[\epsilon_j + \omega_{\lambda i} - R(Jj\lambda i) - \eta_{J\nu}] - \frac{1}{\sqrt{2}}\Gamma(Jj\lambda i)\} \times$$

$$\times[1 + \mathcal{L}(Jj\lambda i)] + \sum_{12} F_{j\lambda}^{\lambda_1 i_1\lambda_2 i_2}(J\nu)U_{\lambda_1 i_1}^{\lambda_2 i_2}(\lambda i) \times$$

$$\times[1 + \frac{1}{2}K^{\lambda}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)][1 + \mathcal{L}(Jj\lambda i)] -$$

$$-\sqrt{2}\sum_{13} F_{j_3\lambda_3}^{\lambda_1 i_1\lambda i}(J\nu)\hat{j}\hat{\lambda}_3(-)^{j_3+\lambda_1+\lambda+J}\begin{Bmatrix}\lambda_1 & \lambda & \lambda_3\\ J & j_3 & j\end{Bmatrix}\Gamma(jj_3\lambda_1 i_1) \times$$

$$\times[1 + \frac{1}{2}K^{\lambda_3}(\lambda i, \lambda_1 i_1|\lambda_1 i_1, \lambda i)][1 + \mathcal{L}(Jj_3\lambda_1 i_1) + \mathcal{L}(Jj\lambda i)] = 0, \qquad (17)$$

$$F_{j\lambda}^{\lambda_1 i_1\lambda_2 i_2}(J\nu)[\epsilon_j + \omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2} - S(Jj\lambda|\lambda_1 i_1, \lambda_2 i_2) - \eta_{J\nu}] \times$$

$$\times[1 + \frac{1}{2}K^{\lambda}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)][1 + \mathcal{M}(Jj\lambda|\lambda_1 i_1, \lambda_2 i_2)] +$$



$$+ \frac{1}{2}\sum_i D_j^{\lambda i}(J\nu)U_{\lambda_1 i_1}^{\lambda_2 i_2}(\lambda i)[1 + \frac{1}{2}K^{\lambda}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)] \times$$

$$\times[1 + \mathcal{L}(Jj\lambda i)] - \frac{1}{2\sqrt{2}}\sum_{i_2}\{D_{j_2}^{\lambda_2 i_2}(J\nu)\Gamma(j_2 j\lambda_1 i_1)\hat{j}_2\hat{\lambda} \times$$

$$\times(-)^{j+\lambda_1+\lambda_2+J}\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda\\ J & j & j_2\end{Bmatrix}[1 + \frac{1}{2}K^{\lambda}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)] \times$$

$$\times[1 + \mathcal{L}(Jj\lambda_1 i_1) + \mathcal{L}(Jj_2\lambda_2 i_2)] + D_{j_2}^{\lambda_1 i_1}(J\nu)\Gamma(j_2 j\lambda_2 i_2)\hat{j}_2\hat{\lambda} \times$$

$$\times(-)^{j+J-\lambda}\begin{Bmatrix}\lambda_2 & \lambda_1 & \lambda\\ J & j & j_2\end{Bmatrix}[1 + \frac{1}{2}K^{\lambda}(\lambda_2 i_2, \lambda_1 i_1|\lambda_1 i_1, \lambda_2 i_2)] \times$$

$$\times[1 + \mathcal{L}(Jj\lambda_2 i_2) + \mathcal{L}(Jj_2\lambda_1 i_1)]\} = 0. \qquad (18)$$

We obtain the basic system of QPM equations for odd spherical nuclei, substituting the expression for $F_{j\lambda}^{\lambda_1 i_1\lambda_2 i_2}(J\nu)$ from (18) into (16), (17).

$$\mathfrak{F}(\eta_{J\nu}) \equiv \epsilon_J - \eta_{J\nu} - \frac{1}{2}\sum_{\lambda ij} D_j^{\lambda i}(J\nu)\Gamma(Jj\lambda i)[1 + \mathcal{L}(Jj\lambda i)] = 0, \qquad (19)$$

$$D_{j_1}^{\lambda_1 i_1}(J\nu)[\epsilon_{j_1} + \omega_{\lambda_1 i_1} - R(Jj_1\lambda_1 i_1) - \eta_{J\nu}][1 + \mathcal{L}(Jj_1\lambda_1 i_1)] -$$

$$-\frac{1}{2}\sum_{23}\frac{(2\lambda_3+1)\hat{j}_1\hat{j}_3\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda_3\\ j_2 & J & j_1\end{Bmatrix}\Gamma(j_1 j_2\lambda_2 i_2)}{[\epsilon_{j_2} + \omega_{\lambda_2 i_2} + \omega_{\lambda_1 i_1} - S(Jj_2\lambda_3|\lambda_2 i_2, \lambda_1 i_1) - \eta_{J\nu}]} \times$$

$$\times\frac{[1 + \frac{1}{2}K^{\lambda_3}(\lambda_1 i_1, \lambda_2 i_2|\lambda_2 i_2, \lambda_1 i_1)][1 + \mathcal{L}(Jj_2\lambda_2 i_2) + \mathcal{L}(Jj_1\lambda_1 i_1)]}{[1 + \mathcal{M}(Jj_2\lambda_3|\lambda_2 i_2, \lambda_1 i_1)]} \times$$

$$\times\{D_{j_3}^{\lambda_1 i_1}(J\nu)\Gamma(j_3 j_2\lambda_2 i_2)\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda_3\\ j_1 & J & j_3\end{Bmatrix}[1 + \mathcal{L}(Jj_2\lambda_2 i_2) + \mathcal{L}(Jj_3\lambda_1 i_1)] +$$

$$+ D_{j_3}^{\lambda_2 i_2}(J\nu)\Gamma(j_3 j_2\lambda_1 i_1)\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda_3\\ J & j_2 & j_3\end{Bmatrix}(-)^{\lambda_1+\lambda_2-\lambda_3}[1 + \mathcal{L}(Jj_2\lambda_1 i_1) + \mathcal{L}(Jj_3\lambda_2 i_2)]\}$$

$$-\frac{1}{2}\sum_{234}\frac{D^{\lambda_1 i_4}_{j_1}(J\nu)\,U^{\lambda_3 i_3}_{\lambda_2 i_2}(\lambda_1 i_1)\,U^{\lambda_3 i_3}_{\lambda_2 i_2}(\lambda_1 i_4)}{[\epsilon_{j_1}+\omega_{\lambda_2 i_2}+\omega_{\lambda_3 i_3}-S(Jj_1\lambda_1|\lambda_2 i_2,\lambda_3 i_3)-\eta_{J\nu}]}\times$$

$$\times\frac{[1+\frac{1}{2}K^{\lambda_1}(\lambda_3 i_3,\lambda_2 i_2|\lambda_2 i_2,\lambda_3 i_3)][1+\mathfrak{L}(Jj_1\lambda_1 i_4)][1+\mathfrak{L}(Jj_1\lambda_1 i_1)]}{[1+\mathfrak{M}(Jj_1\lambda_1|\lambda_2 i_2,\lambda_3 i_3)]}+$$

$$+\frac{1}{\sqrt{2}}\sum_{23} D^{\lambda_3 i_3}_{j_3}(J\nu)\,\Gamma(j_3 j_1\lambda_2 i_2)\,\hat{j}_3(-)^{j_1+\lambda_1+J}\begin{Bmatrix}\lambda_1 & \lambda_2 & \lambda_3\\ j_3 & J & j_1\end{Bmatrix}\times$$

$$\times\Big\{\frac{\hat{\lambda}_1\,U^{\lambda_2 i_2}_{\lambda_3 i_3}(\lambda_1 i_1)[1+\frac{1}{2}K^{\lambda_1}(\lambda_3 i_3,\lambda_2 i_2|\lambda_2 i_2,\lambda_3 i_3)][1+\mathfrak{L}(Jj_1\lambda_1 i_1)]}{[1+\mathfrak{M}(Jj_1\lambda_1|\lambda_2 i_2,\lambda_3 i_3)]}\times$$

$$\times\frac{[1+\mathfrak{L}(Jj_1\lambda_2 i_2)+\mathfrak{L}(Jj_3\lambda_3 i_3)]}{[\epsilon_{j_1}+\omega_{\lambda_2 i_2}+\omega_{\lambda_3 i_3}-S(Jj_1\lambda_1|\lambda_2 i_2,\lambda_3 i_3)-\eta_{J\nu}]}+$$

$$+\frac{\hat{\lambda}_3\,U^{\lambda_1 i_1}_{\lambda_2 i_2}(\lambda_3 i_3)[1+\frac{1}{2}K^{\lambda_3}(\lambda_1 i_1,\lambda_2 i_2|\lambda_2 i_2,\lambda_1 i_1)][1+\mathfrak{L}(Jj_3\lambda_3 i_3)]}{[\epsilon_{j_3}+\omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2}-S(Jj_3\lambda_3|\lambda_2 i_2,\lambda_1 i_1)-\eta_{J\nu}]}\times$$

$$\times\frac{[1+\mathfrak{L}(Jj_3\lambda_2 i_2)+\mathfrak{L}(Jj_1\lambda_1 i_1)]}{[1+\mathfrak{M}(Jj_3\lambda_3|\lambda_2 i_2,\lambda_1 i_1)]}\Big\}=\frac{1}{\sqrt{2}}\Gamma(Jj_1\lambda_1 i_1)[1+\mathfrak{L}(Jj_1\lambda_1 i_1)]. \qquad (20)$$

With coefficients $D^{\lambda i}_{j}$ found from the system of equations (20), one can solve (19) and find the energy $\eta_{J\nu}$ for states $\Psi_\nu$ (JM).

It should be noted that equation (19) has the same form even with taking into account more complex components in the wave function (1). Actually, this equation is the matrix form of Dyson equation in the QPM. One could obtain a similar system of equations in nondiagonal approximation using (7) and (9).

If we assume that the odd quasiparticle does not influence the structure of two-phonon excitations of the even-even core and neglect Pauli principle corrections, the $\mathfrak{L}, K^\lambda$ and $\mathfrak{M}$ functions become zero, and our system of equations is reduced to the system of equations used in [4] for calculation of radiative strength functions in odd spherical nuclei. Further, if we drop



all terms containing $U^{\lambda_2 i_2}_{\lambda_3 i_3}(\lambda i)$, which are important for the description of fragmentation of "quasiparticle plus phonon" states [4], we get the system of equations firstly obtained in [8]. Neglecting of terms containing U corresponds to the usual assumptions of phenomenological models [11]. Such equations have been used for calculation of deep hole strength distributions and neutron strength functions [3,4]. If "quasiparticle plus two phonon" components are not included into the wave function of an excited state of an odd nucleus, all terms of (20), except those in the first line, disappear, and one gets the system of equations firstly obtained in [9]. Then, assuming $\mathfrak{L}(Jj\lambda i)=0$ one gets the well-known secular equation of the superfluidity nuclear model [12]. Comparison of equations (19), (20) with those from [3,4,8,9,12] shows that the inclusion of Pauli principle leads to the renormalization of matrix elements and to the shifts in the energetic poles. In the number of cases, these corrections play an important role [5,7,8].

Thus, our system of QPM equations is most general and it contains all approximate equations used in the calculations for odd spherical nuclei within the QPM [3,4,9]. These equations may be used for calculations of various characteristics of low- and high-lying states in odd spherical nuclei.

Table

$\mathfrak{M}^J_2(j_1\lambda'_1,\lambda_1 i_1|\lambda_2 i_2, j_2\lambda'_2)$ values for states with $J^\pi=9/2^+$ in $^{209}$Pb and $J^\pi=1/2^-$ in $^{61}$Ni

| Nucleus | $j_1$ | $\lambda'_1$ | $\lambda_1 i_1$ | $\lambda_2 i_2$ | $j_2$ | $\lambda'_2$ | $\mathfrak{M}^J_2$ |
|---|---|---|---|---|---|---|---|
| $^{209}$Pb | $1i_{13/2}$ | 2 | $2^+_1$ | $2^+_1$ | $1i_{13/2}$ | 2 | -0.018 |
| | $1i_{13/2}$ | 2 | $2^+_1$ | $2^+_2$ | $1i_{13/2}$ | 2 | 0.006 |
| | $1i_{13/2}$ | 2 | $3^-_1$ | $2^+_3$ | $1i_{13/2}$ | 2 | -0.0001 |
| | $1i_{13/2}$ | 2 | $2^+_2$ | $2^+_2$ | $1i_{13/2}$ | 2 | 0.003 |
| $^{61}$Ni | $2p_{3/2}$ | 2 | $2^+_1$ | $2^+_2$ | $2p_{3/2}$ | 2 | -0.67 |
| | $2p_{3/2}$ | 2 | $2^+_1$ | $2^+_3$ | $2p_{3/2}$ | 2 | 0.12 |
| | $2p_{3/2}$ | 2 | $2^+_2$ | $2^+_2$ | $2p_{3/2}$ | 2 | 1.12 |
| | $2p_{3/2}$ | 2 | $2^+_2$ | $2^+_4$ | $2p_{3/2}$ | 2 | 0.18 |

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Дао Тиен Кхоа, Вдовин А.И., Воронов В.В. E4-84-501
Основные уравнения квазичастично-фононной модели ядра для нечетных сферических ядер

Получена в общем виде система основных уравнений квазичастично-фононной модели для нечетных сферических ядер. Уравнения учитывают эффекты ангармоничности колебаний четно-четного остова и поправки, обусловленные требованием выполнения принципа Паули. Показано, что выведенная система уравнений содержит все варианты приближенных уравнений квазичастично-фононной модели, широко используемые в расчетах.

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