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DESCRIPTION OF LOW-LYING STATES IN ODD-ODD DEFORMED NUCLEI TAKING ACCOUNT OF THE COUPLING WITH CORE ROTATIONS AND VIBRATIONS

I. Theory

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1. Introduction

Low-lying states in odd-odd nuclei have been extensively investigated for a long time, both empirically and theoretically [1-22]. The main attention has been paid to the Gallagher-Moszkowski splitting [1] and Newby shift [2] which are directly connected with neutron-proton interaction and provide a good possibility for its investigation. In most of the papers the neutron-proton interaction was introduced as some effective forces with parameters fitted in accordance with the available experimental data (see, for example, [6]). In [6,10-14] the Coriolis mixing was involved that improved considerably the description of low-lying states in deformed odd-odd nuclei. Recently, a successful attempt has been made to derive the microscopic description of odd-odd nuclei within the rotor-plus-two-quasiparticle approach, where the single-particle states, static equilibrium core properties, and the residual neutron-proton interaction were determined from the same nucleon-nucleon interaction [9].

As a rule, the coupling of external nucleons with even-even core vibrations was neglected. On the other hand, this coupling leads to the appearance in low-lying states of odd-odd nuclei of vibrational admixtures [4], which may be very important, especially, for the description of $E\lambda(M\lambda)$-transitions. A similar effect was clearly demonstrated in numerous calculations within the QPM [25-28] of $E\lambda$-transition rates in odd deformed nuclei [27-30]. The importance of vibrational admixtures in the states of odd-odd nuclei has been also confirmed in some phenomenological models [18-21]. It should be noted that just vibrational admixtures are mainly responsible for the well known attenuation of Coriolis interaction matrix elements [29,30].

The up-to-day status of investigations of low-lying states in deformed odd-odd nuclei clearly requires the construction of a general microscopic approach, which would include consistently the neutron-proton interaction and the coupling with rotational and vibrational core excitations in the framework of a common microscopic scheme and, on the other hand, would be able to describe the properties of even-even, odd-A and odd-odd nuclei on the same microscopic footing. Up to now, such a microscopic approach is absent but it can be derived on the basis of the QPM [4] in the same way as for odd nuclei [29]. The QPM seems to be the most suitable for this aim since just this model has been successfully applied for the description of low-lying states in a wide region of even-even and odd nuclei [25-29,31]. In [29,30] the Coriolis interaction was included into the QPM. Some cases of interest were brought to the light. In particular, it has been shown that in odd Eu-Tb isotopes only the simultaneous use of Coriolis and quasiparticle-phonon interactions enables us to describe the anomalous behaviour of E1-transitions [29].

The aim of this paper is the formulation, on the QPM basis [4], of the general microscopic approach for the description of low-lying states in odd-odd deformed nuclei. This approach will include the coupling of external nucleons with even-even core vibrations due to the quasiparticle-phonon interaction, the rotational excitations as well as the Coriolis mixing, and the interaction between external proton and neutron, which results in the Gallagher-Moszkowski splitting and the Newby shift. It should be noted that the latter appears in our approach as a part of residual forces, omitted earlier in the QPM [4]. In addition to the QPM version [4], we take into account the mixing of neutron-proton configurations due to the quasiparticle-phonon interaction. The particle-particle (pairing) and isovector parts of this interaction are also included.

In Sec. 2, the rotational part of the approach is outlined. In Sec. 3, the intrinsic Hamiltonian is considered and the corresponding part of neutron-proton interaction leading to the Gallagher-Moszkowski splitting and Newby shift is extracted. In Sec. 4, the secular equation for excitation energies and the expressions for the wave function coefficients of nonrotational states are derived. In sec. 5, the Gallagher-Moszkowski splitting and Newby shift are considered. The expressions for $E\lambda(M\lambda)$ transition rates are presented in Sec. 6. A short discussion and main conclusions are given in Sec. 7. In Appendices A and B the matrix elements of the rotational Hamiltonian and some expressions for intrinsic excitations are presented, respectively.

2. The rotational excitations

The Hamiltonian of our approach is written as a sum of intrinsic and rotational parts:

$$H = H_{intr} + H_{rot} . \tag{1}$$

The H_{intr} will be considered in Sec. 3. The rotational part of the Hamiltonian (1) includes the rotation of the nucleus as a whole, the Coriolis interaction and the centrifugal term:

$$H_{rot} = H_R + H_{CI} + H_j , \tag{2}$$

where

$$H_R = \frac{\hbar^2}{2\Phi} (\hat{I}^2 - \hat{I}_3^2), \tag{3.1}$$

$$H_{CI} = -\frac{\hbar^2}{2\Phi} (I^+ j^- + I^- j^+), \tag{3.2}$$

$$H_j = \frac{\hbar^2}{2\Phi} (j^+ j^- + j^- j^+). \tag{3.3}$$

In (3) Φ is moment of inertia of the odd-odd nucleus; $\hat{I}_3$ and $\hat{j}_3$ are operators of projection of the total ($\vec{I}$) and intrinsic ($\vec{j} = \vec{j}_n + \vec{j}_p$) angular momentum into nuclear symmetry axis, respectively; $I^{\pm}$ and $j^{\pm}$ the corresponding momentum shift operators.

The wave function of the state of odd-odd nucleus has the form

$$|I^{\pi} M \rho> = \sum_{K\nu} b_{\nu K}^{I\rho} |I^{\pi} M K \nu>, \tag{4}$$

where $b_{\nu K}^{I\rho}$ are the Coriolis mixing coefficients; M and K are the angular momentum projections in laboratory and intrinsic systems, respectively; ρ and ν are the additional quantum numbers. Further [32],

$$|I^{\pi} M K \nu> = \sqrt{\frac{2I+1}{16\pi^2(1+\delta_{K,0})}} \left(\mathcal{D}_{MK}^{I} + (-1)^{I+K} \mathcal{D}_{M-K}^{I} R_i\right) \Psi_{\nu}(K^{\pi}), \tag{5}$$

where $\Psi_{\nu}(K^{\pi})$ is the eigenvector of H_{intr}; R_i is the operator of rotation by angle π around the second intrinsic axis.

The matrix elements of different parts of the Hamiltonian between the states (5) are presented in Appendix A. We don't give the derivation of them as it can be found in [6,9]. Note that instead of single-particle matrix elements $\langle n'|j^+_n|n\rangle$ and $\langle p'|j^+_p|p\rangle$, as in [6,9], the expressions in Appendix A contain the matrix elements $\langle\Psi_{\nu'}(K'^{\pi})|j^+|\Psi_\nu(K^\pi)\rangle$ where the intrinsic wave function $\Psi_\nu(K^\pi)$ includes two-quasiparticle (neutron-proton) as well as two-quasiparticle $\otimes$ phonon components. The latter are of vibrational type. They lead to decreasing of the amplitude of main two-quasiparticle component in $\Psi_\nu(K^\pi)$ and, as a result, to attenuation of the Coriolis matrix elements, observed experimentally [29,30]. Thus, in the framework of our approach the attenuation effect is described in a natural way on the microscopic footing.

3. The intrinsic Hamiltonian and neutron-proton interaction

In accordance with [25,26,33] the intrinsic part of the Hamiltonian (1) is written as

$$H_{intr} = H_{sp} + H_{pair} + H_{mm}, \tag{6}$$

where

$$H_{sp} = \sum_\tau \sum_{\hat{q}\in\tau} (E_q - \lambda_\tau)\, a^+_{\hat{q}} a_{\hat{q}} \tag{7}$$

is a single-particle potential,

$$H_{pair} = -\sum_\tau G_\tau \sum_{qq'\in\tau} a^+_{q+} a^+_{q-} a_{q'-} a_{q'+} \tag{8}$$

is the monopole pairing and

$$H_{mm} = -1/2 \sum_{\lambda\hat{\mu}} \sum_{\tau\tau'} (\varkappa_0^{(\lambda\mu)} + \tau\tau'\varkappa_1^{(\lambda\mu)})\, Q^{(\tau)}_{\lambda\hat{\mu}} Q^{(\tau')}_{\lambda-\hat{\mu}} \tag{9}$$

is the multipole isoscalar and isovector interaction with multipole operator

$$Q^{(\tau)}_{\lambda\hat{\mu}} = \sum_{\hat{q}_1\hat{q}_2\in\tau} \delta_{\hat{K}_1-\hat{K}_2,\hat{\mu}} \langle\hat{q}_1|\hat{f}^{\lambda\mu}|\hat{q}_2\rangle\, a^+_{\hat{q}_1} a_{\hat{q}_2}. \tag{10}$$

In (7)-(10) we used the following notation: τ means neutron and proton systems for which τ = -1 and +1, respectively; $a^+_{\hat{q}}$ is the particle creation operator for single-particle state $\hat{q}$; $\hat{q} = q\sigma$, $\hat{K} = K\sigma$, $\hat{\mu} = \mu\sigma$, $K \geq 0$, $\mu \geq 0$; $\sigma = \pm 1$ characterises the symmetry with respect to time reversal operation; E_q is single-particle energy; G_τ and λ_τ are pairing strength constant and chemical potential; $\varkappa_0^{(\lambda\mu)}$ and $\varkappa_1^{(\lambda\mu)}$ are the multipole isoscalar and isovector strength constants, respectively;

$\langle \hat{q}_1 | \hat{f}^{\lambda\mu} | \hat{q}_2 \rangle$ is the single-particle matrix element for the operator

$$\hat{f}^{\lambda\mu} = R(r)\{ Y_{\lambda\mu} + (-1)^{\mu} Y_{\lambda-\mu} \} (1+\delta_{\mu,0})^{-1} \tag{11}$$

with unspecified radial dependence $R(r)$.

After the Bogoliubov transformation and using the RPA equations for one-phonon excitations of even-even core, the intrinsic Hamiltonian (6) can be transformed to the form [25,26,33]:

$$H_{intr} = H_{\alpha+Q} + H_{QB} + H^{pair}_{QB} + H_{BB}\,, \tag{12}$$

where

$$H_{\alpha+Q} = \sum_q \mathcal{E}_q B(qq0) - \frac{1}{4} \sum_{\lambda\hat{\mu}} \sum_{ii'} \sum_{\tau} \frac{X^g_\tau + X^{g'}_\tau}{\sqrt{\mathcal{Y}^g_\tau \mathcal{Y}^{g'}_\tau}} Q^+_{\hat{g}} Q_{\hat{g}} \tag{13.1}$$

generates quasiparticle excitations and phonon excitations of a doubly-even core,

$$H_{QB} = -\frac{1}{4} \sum_{\hat{g}} \Big\{ (Q^+_{\hat{g}} + Q_{-\hat{g}}) \sum_{\tau} \sum_{q_1 q_2 \in \tau} \Gamma^g_{q_1 q_2} B(q_1 q_2 -\hat{\mu}) + h.c. \Big\} \tag{13.2}$$

is the quasiparticle-phonon interaction which will be shown to mix neutron-proton and neutron-proton ⊗ phonon configurations in the wave function of odd-odd nucleus,

$$H^{pair}_{QB} = \frac{1}{\sqrt{2}} \sum_{\tau} G_\tau \sum_{q_1 q_2 \in \tau} u_{q_1} v_{q_1} (u^2_{q_2} - v^2_{q_2}) \sum_i \Big\{ (\psi^{20i}_{q_2 q_2} Q^+_{20i} + \varphi^{20i}_{q_2 q_2} Q_{20i}) B(q_1 q_1 0) + h.c. \Big\} \tag{13.3}$$

is the pairing quasiparticle-phonon interaction and

$$H_{BB} = -\frac{1}{2} \sum_{\lambda\hat{\mu}} \sum_{\tau\tau'} (\varkappa_0^{(\lambda\mu)} + \tau\tau' \varkappa_1^{(\lambda\mu)}) \sum_{\substack{q_1 q_2 \in \tau \\ q_1' q_2' \in \tau}} f^{\lambda\mu}_{q_1 q_2} f^{\lambda\mu}_{q_1' q_2'} v_{q_1 q_2} v_{q_1' q_2'} B(q_1 q_2 \hat{\mu}) B(q_1' q_2' -\hat{\mu}) \tag{13.4}$$

is the interaction which will be shown to be responsible for the Gallagher-Moszkowski splitting and Newby shift. In (13.1)-(13.4) the following notation is used:

$$Q^+_{\hat{g}} = \frac{1}{2} \sum_{q_1 q_2} \{ \psi^g_{q_1 q_2} A^+(q_1 q_2 \hat{\mu}) - \varphi^g_{q_1 q_2} A(q_1 q_2 -\hat{\mu}) \} \tag{14}$$

is the creation operator of one-phonon state $\hat{g} \equiv g\sigma \equiv \lambda\mu i\sigma$, where i is the number of the phonon with given $\lambda\mu$ and

$$A^{+}(q_1 q_2 \hat{\mu}) = -\frac{1}{\sqrt{1+\delta_{\mu,0}}} \sum_{\sigma_1 \sigma_2} \delta_{\hat{K}_1+\hat{K}_2,\hat{\mu}} \, \alpha^{+}_{\hat{q}_1} \alpha^{+}_{\hat{q}_2} \, \theta_{\sigma_1,-\sigma_2} , \tag{15.1}$$

$$B(q_1 q_2 \hat{\mu}) = \sum_{\sigma_1 \sigma_2} \delta_{\hat{K}_1+\hat{K}_2,\hat{\mu}} \, \alpha^{+}_{\hat{q}_1} \alpha_{-\hat{q}_2} \, \theta_{-\sigma_1,-\sigma_2} . \tag{15.2}$$

Here, $\alpha^{+}_{\hat{q}}$ is the quasiparticle creation operator; $u_{q_1 q_2} = u_{q_1} v_{q_2} + v_{q_1} u_{q_2}$, $v_{q_1 q_2} = u_{q_1} u_{q_2} - v_{q_1} v_{q_2}$, where u_q and v_q are the Bogoliubov transformation coefficients; $f^{\lambda\mu}_{q_1 q_2}$ is the single-particle matrix elements of operator (11); $\theta_{\sigma_1 \sigma_2} = 1 - 2\delta_{\sigma_1,1}\delta_{\sigma_2,1}$. The expressions for the functions X^{g}_{τ}, Y^{g}_{τ} and $\Gamma^{g}_{q_1 q_2}$ are given in Appendix B.

It is easy to see that if we consider the general case of neutron-proton interaction V_{np} which is proposed to be hermitian and invariant under time reversal

$$V^{tot}_{np} = \frac{1}{2} \sum_{\substack{\hat{r}\hat{r}' \\ \hat{s}\hat{s}'}} \langle \hat{r}\hat{s} | \hat{V}_{np} | \hat{r}'\hat{s}' \rangle \, a^{+}_{\hat{r}} a_{\hat{r}'} a^{+}_{\hat{s}} a_{\hat{s}'} , \tag{16}$$

then after the Bogoliubov transformation this interaction can be expressed as a sum of $\alpha^{+}\alpha$, $\alpha^{+}\alpha^{+}$, $\alpha^{+}\alpha^{+}\alpha\alpha$, $\alpha^{+}\alpha^{+}\alpha^{+}\alpha$ and $\alpha^{+}\alpha\alpha^{+}\alpha$ type terms and of their h.c. counterparts (in (16) labels r and s mean proton and neutron particle (quasiparticle) states, respectively). It is clear that all the terms of the sum, except $\alpha^{+}\alpha\alpha^{+}\alpha$ type term having the form

$$V_{np} = \frac{1}{2} \sum_{\substack{\hat{r}\hat{r}' \\ \hat{s}\hat{s}'}} \langle \hat{r}\hat{s} | \hat{V}_{np} | \hat{r}'\hat{s}' \rangle \Big\{ (u_r u_{r'} u_s u_{s'} + v_r v_{r'} v_s v_{s'}) \alpha^{+}_{\hat{r}} \alpha_{\hat{r}'} \alpha^{+}_{\hat{s}} \alpha_{\hat{s}'} - (v_r v_{r'} u_s u_{s'} + u_r u_{r'} v_s v_{s'}) \sigma_r \sigma_{r'} \alpha^{+}_{-\hat{r}} \alpha_{-\hat{r}'} \alpha^{+}_{\hat{s}} \alpha_{\hat{s}'} \Big\} , \tag{17}$$

will contribute to (13.1)-(13.3) parts of the Hamiltonian (12), while the term (17), the diagonal ($r = r'$, $s = s'$) matrix elements of which are used to describe the Gallagher-Moszkowski splitting and the Newby shift (see, for example, [9]), will contribute to the n-p part of the H_{BB} (13.4). The n-p part of the H_{BB} and n-p interaction (17) have the same quasiparticle structure and the same physical origin. Therefore, the neutron-proton interaction of interest is not introduced in our approach from outside, but appears as the inherent part of the microscopic Hamiltonian (6). It should be noted that the interaction H_{QQ} (as well as the interaction H_{QQ}^{pair}) is used to be omitted in the QPM calculations. For the interaction (17) is more general than the n-p part of the H_{QQ} (the former may contain different types of residual interaction, while the latter is written only for separable multipole forces), we will use V_{np} (17) in the following consideration instead of the H_{QQ}.

4. The main equations for the intrinsic excitations

The intrinsic wave function is given as

$$\Psi_{\nu\gamma_0}(\tilde{K}_0^\pi)=\Big\{\sum_{\tilde{s}\tilde{r}} C_{sr}^{\nu\gamma_0}\hat{A}^+_{\gamma_0}(\tilde{s}\tilde{r}\tilde{K}_0)+\sum_{\tilde{s}\tilde{r}\tilde{g}\gamma} k_\mu^{K_0} D_{srg}^{\nu\gamma_0}\hat{A}^+_\gamma(\tilde{s}\tilde{r}\tilde{K})Q^+_{\tilde{g}}\,\delta_{\tilde{K}+\tilde{\mu},\tilde{K}_0}\Big\}|\rangle, \tag{18}$$

where

$$\tilde{A}^+_\gamma(\tilde{s}\tilde{r}\tilde{K})=\frac{1}{\sqrt{1+\delta_{K,0}}}\alpha^+_{\tilde{s}}\alpha^+_{\tilde{r}}\,\delta_{\tilde{K}_s+\tilde{K}_r,\tilde{K}}\left(1-(1+\gamma)\delta_{K,0}\,\delta_{\sigma_s,-1}\delta_{\sigma_s,+1}\right), \tag{19.1}$$

$$k_\mu^{K_0}=\left(1+\delta_{K_0,0}(1-\delta_{\mu,0})\right)^{-1/2}, \tag{19.2}$$

$|\rangle$ is the vacuum for quasiparticle and phonon operators ($\alpha_{\hat{q}}|\rangle=Q_{\hat{g}}|\rangle=0$), ν is the number of intrinsic state with given K_0^π. Note that for $K_0=0$ the function $\Psi_{\nu\gamma_0}(\tilde{K}_0^\pi)$ includes the eigenvalue $\gamma_0=\pm 1$ of the operator R_i and fulfils the condition (A2) (see Appendix A).

Let us consider the selection conditions $|K_s\pm K_r|=K_0$ and $|K\pm\mu|=K_0$ to be embedded into the amplitudes $C_{sr}^{\nu\gamma_0}$ and $D_{srg}^{\nu\gamma_0}$, respectively. Then, the normalization condition for the wave function (18) is

$$\left(\Psi^*_{\nu\gamma_0}(\tilde{K}_0^\pi)\Psi_{\nu\gamma_0}(\tilde{K}_0^\pi)\right)=\sum_{sr}\left(C_{sr}^{\nu\gamma_0}\right)^2+\sum_{srg}\left(D_{srg}^{\nu\gamma_0}\right)^2=1. \tag{20}$$

The amplitudes $C_{sr}^{\nu\gamma_0}$ and $D_{srg}^{\nu\gamma_0}$ can be obtained by variational method with keeping the condition (20):

$$\delta\Big\{\left(\Psi^*_{\nu\gamma_0}(\tilde{K}_0^\pi)H_{intr}\Psi_{\nu\gamma_0}(\tilde{K}_0^\pi)\right)-\eta_{\nu\gamma_0}\left(\left(\Psi^*_{\nu\gamma_0}(\tilde{K}_0^\pi)\Psi^*_{\nu\gamma_0}(\tilde{K}_0^\pi)\right)-1\right)\Big\}=0, \tag{21}$$

where amplitudes $C_{sr}^{\nu\gamma_0}$ and $D_{srg}^{\nu\gamma_0}$ are the variational variables and Lagrange multiplier $\eta_{\nu\gamma_0}$ means the excitation energy of the state (18).

Using the expectation value of H_{intr}

$$
\begin{aligned}
(\Psi^*_{\nu\gamma_0}(\hat{K}^\pi_0) H_{intr} \Psi_{\nu\gamma_0}(\hat{K}^\pi_0)) = \\
= \sum_{sr} (C^{\nu\gamma_0}_{sr})^2 (\varepsilon_s + \varepsilon_r) + \sum_{srg} (\mathcal{D}^{\nu\gamma_0}_{srg})^2 (\varepsilon_s + \varepsilon_r + \omega_g) - \\
- \sum_{sr} \sum_{s'r'g} C^{\nu\gamma_0}_{sr} \mathcal{D}^{\nu\gamma_0}_{s'r'g} (\hat{\Gamma}^g_{ss'} \delta_{r,r'} + \hat{\Gamma}^g_{rr'} \delta_{s,s'}) + \quad (22) \\
+ \sum_{sr} \sum_{s'r'} C^{\nu\gamma_0}_{sr} C^{\nu\gamma_0}_{s'r'} \langle rs | V_{np} | r's' \rangle_{0\gamma_0} + \sum_{sr} \sum_{\substack{s'r'g \\ \gamma\gamma'}} \mathcal{D}^{\nu\gamma_0}_{srg} \mathcal{D}^{\nu\gamma_0}_{s'r'g} \langle rs | V_{np} | r's' \rangle_{\mu\gamma\gamma'}
\end{aligned}
$$

we will have from (21) the system of equations for the amplitudes $C^{\nu\gamma_0}_{sr}$ and $\mathcal{D}^{\nu\gamma_0}_{srg}$:

$$
\begin{aligned}
2 C^{\nu\gamma_0}_{sr} (\varepsilon_s + \varepsilon_r - \eta_{\nu\gamma_0}) - \sum_{s'r'g} \mathcal{D}^{\nu\gamma_0}_{s'r'g} (\hat{\Gamma}^g_{ss'} \delta_{r,r'} + \hat{\Gamma}^g_{rr'} \delta_{s,s'}) + \quad (23.1) \\
+ \sum_{s'r'} C^{\nu\gamma_0}_{s'r'} (\langle rs | V_{np} | r's' \rangle_{0\gamma_0} + \langle r's' | V_{np} | rs \rangle_{0\gamma_0}) = 0,
\end{aligned}
$$

$$
\begin{aligned}
2 \mathcal{D}^{\nu\gamma_0}_{srg} (\varepsilon_s + \varepsilon_r + \omega_g - \eta_{\nu\gamma_0}) - \sum_{s'r'} C^{\nu\gamma_0}_{s'r'} (\hat{\Gamma}^g_{ss'} \delta_{r,r'} + \hat{\Gamma}^g_{rr'} \delta_{s,s'}) + \quad (23.2) \\
+ \sum_{s'r'\gamma\gamma'} \mathcal{D}^{\nu\gamma_0}_{s'r'g} (\langle rs | V_{np} | r's' \rangle_{\mu\gamma\gamma'} + \langle r's' | V_{np} | rs \rangle_{\mu\gamma\gamma'}) = 0.
\end{aligned}
$$

(the expressions for $\tilde{\Gamma}^g_{qq'}$ and diagonal case of $\langle rs | V_{np} | rs \rangle_{\mu\gamma\gamma'}$ can be found in Appendix B).

If we neglect the nondiagonal matrix elements $\langle rs | V_{np} | r's' \rangle_{\mu\gamma\gamma'}$ of n-p interaction equations (23.1)÷(23.2) are simplified and can be rewritten as

$$
\begin{aligned}
\mathcal{D}^{\nu\gamma_0}_{srg} = (\varepsilon_s + \varepsilon_r + \omega_g + \langle rs | V_{np} | rs \rangle_{\mu\gamma} - \eta_{\nu\gamma_0})^{-1} \times \quad (24) \\
\times \tfrac{1}{2} \sum_{s'r'} C^{\nu\gamma_0}_{s'r'} (\hat{\Gamma}^g_{ss'} \delta_{r,r'} + \hat{\Gamma}^g_{rr'} \delta_{s,s'})
\end{aligned}
$$

$$\sum_{s'r'} C^{\nu\gamma_0}_{s'r'} \left\{ \delta_{sr,s'r'} \left(\mathcal{E}_s + \mathcal{E}_r + \langle rs | V_{np} | rs \rangle_{o\gamma_0} - \eta_{\nu\gamma_0} - \frac{1}{4} \sum_{s_1 r_1 g\gamma} \frac{\widetilde{\Gamma}^{g}_{ss_1rr_1} \cdot \widetilde{\Gamma}^{g}_{s's_1r'r_1}}{\mathcal{E}_{s_1} + \mathcal{E}_{r_1} + \omega_g + \langle r_1 s_1 | V_{np} | r_1 s_1 \rangle_{\mu\gamma} - \eta_{\nu\gamma_0}} \right\} \quad (25)$$

with the secular equation for finding the excitation energies

$$\det \left\| \delta_{sr,s'r'} \left(\mathcal{E}_s + \mathcal{E}_r + \langle rs | V_{np} | rs \rangle_{o\gamma_0} - \eta_{\nu\gamma_0} \right) - \frac{1}{4} \sum_{s_1 r_1 g\gamma} \frac{\widetilde{\Gamma}^{g}_{ss_1rr_1} \cdot \widetilde{\Gamma}^{g}_{s's_1r'r_1}}{\mathcal{E}_{s_1} + \mathcal{E}_{r_1} + \omega_g + \langle r_1 s_1 | V_{np} | r_1 s_1 \rangle_{\mu\gamma} - \eta_{\nu\gamma_0}} \right\| = 0, \quad (26)$$

where $\widetilde{\Gamma}^{g}_{ss_1rr_1} = \widetilde{\Gamma}^{g}_{ss_1} \delta_{r,r_1} + \widetilde{\Gamma}^{g}_{rr_1} \delta_{s,s_1}$.

Expressions (24)-(26) give more general description of intrinsic states of odd-odd nuclei than the previous variant of the QPM in [4]. In addition to [4] they include the interaction V_{np}, the isovector as well as the isoscalar parts of multipole forces (9), the particle-particle contribution (see the second term in (B7)) to quasiparticle-phonon interaction, the mixing of n-p configurations due to the quasiparticle-phonon interaction.

In final expressions (24)-(26), the nondiagonal matrix elements of n-p interaction are neglected for the sake of simplicity. This approximation, in spite of its wide application, can not be considered as substantiated. Indeed, values of nondiagonal matrix elements can be of the same order of magnitude as diagonal ones. It is easy to see also from (23.1)-(23.2) that nondiagonal n-p matrix elements result in an additional mixing of n-p configurations in odd-odd nuclei. So, the role of nondiagonal n-p matrix elements needs further careful investigation.

5. Gallagher-Moszkowski splitting and Newby shift

It is easy to show that the Gallagher-Moszkowski splitting and the Newby shift are embodied in the equations for the intrinsic excitations, derived in Sec. 4. To be sure of this, let us consider the simple case when the long-range residual interaction, except for V_{np}, is neglected ($\widetilde{\Gamma}^{g}_{ss_1rr_1} = 0$) . Then, from (17), (B10) and (26) one can write

$$\eta_{\nu\gamma_0} = \mathcal{E}_s + \mathcal{E}_r + \langle rs | V_{np} | rs \rangle_{o\gamma_0} \overset{\mathcal{E}_s + \mathcal{E}_r}{=} - \delta_{K_0,0} \cdot \gamma_0 \langle r{+}s{-} | V_{np} | r{-}s{+} \rangle +$$
$$+ \delta_{K_s+K_r,K_0} \left\{ \langle r{+}s{+} | V_{np} | r{+}s{+} \rangle (U_r^2 U_s^2 + v_r^2 v_s^2) - \langle r{+}s{-} | V_{np} | r{+}s{-} \rangle (v_r^2 U_s^2 + U_r^2 v_s^2) \right\} -$$
$$- \delta_{|K_s-K_r|,K_0} \left\{ \langle r{+}s{-} | V_{np} | r{+}s{-} \rangle (U_r^2 U_s^2 + v_r^2 v_s^2) - \langle r{+}s{+} | V_{np} | r{+}s{+} \rangle (v_r^2 U_s^2 + U_r^2 v_s^2) \right\} \quad (27)$$

Keeping the condition $\gamma_o = (-1)^I$ (see Appendix A) we finally obtain the well-known expression for the Gallagher-Moszkowski splitting energy, corresponding to the case of independent quasiparticles [6,9]:

$$\Delta E = \eta_{\nu\gamma_o} \cdot \delta_{|K_s - K_r|, K_o} - \eta_{\nu\gamma_o} \delta_{K_s + K_r, K_o} = \\ = \langle r+s-|V_{np}|r+s-\rangle - \langle r+s+|V_{np}|r+s+\rangle + \delta_{K_o,0}(-1)^{I+1} \langle r+s-|V_{np}|r-s+\rangle. \tag{28}$$

The last term in (28) is responsible for the Newby shift:

$$\Delta E_{K_o=0} = \eta_{\nu\gamma_o=-1} - \eta_{\nu\gamma_o=+1} = 2\langle r+s-|V_{np}|r-s+\rangle. \tag{29}$$

6. <u>E(M)λ transitions</u>

The reduced transition probabilities for electric $(X=E)$ and magnetic $(X=M)$ transitions of multipolarity λ between the states $|I^\pi M\rho\rangle$ and $|I'^{\pi'} M'\rho'\rangle$ described by the wave functions (4) is [32]

$$B(X\lambda, I^\pi\rho \to I'^{\pi'}\rho') = \\ = \Big| \sum_{\substack{K_o\nu\gamma_o \\ K_o'\nu'\gamma_o'}} b^{I\rho}_{\nu K_o} b^{I'\rho'}_{\nu' K_o'} \Big\{ (IK_o\lambda K_o'-K_o|I'K_o') \langle \Psi_{\nu'\gamma_o'}(K_o'^{\pi'})| \mathcal{M}'(X\lambda, \mu = K_o'-K_o)|\Psi_{\nu\gamma_o}(K_o^{\pi})\rangle + \\ + (-1)^{I+K_o}(I-K_o\lambda K_o'+K_o|I'K_o')\langle \Psi_{\nu'}(K_o'^{\pi'})|\mathcal{M}'(X\lambda,\mu=K_o'+K_o)|\Psi_\nu(K_o^\pi)\rangle \Big\} \Big|^2. \tag{30}$$

In the intrinsic system, the operator of $X\lambda$ - transition can be written in the case $X=E$ as [33]

$$\mathcal{M}'(E\lambda,\hat{\mu}) = 2\sum_q p^{\lambda\mu}_{qq}(E) v_q^2 + \sum_i L_g^{(E)}(Q^+_{\hat{g}} + Q_{-\hat{g}}) + \\ + \sum_{q_1q_2} p^{\lambda\mu}_{q_1q_2}(E)(u_{q_1}u_{q_2} - v_{q_1}v_{q_2}) B(q_1q_2\hat{\mu}), \tag{31}$$

where

$$L_g^{(E)} = \frac{\sqrt{1+\delta_{\mu,0}}}{2} \sum_{q_1q_2} p^{\lambda\mu}_{q_1q_2}(E)(u_{q_1}v_{q_2} + v_{q_1}u_{q_2})(\Psi^g_{q_1q_2} + \varphi^g_{q_1q_2}) \tag{32}$$

and in the case $X = M$ as

$$\mathcal{M}'(M\lambda,\tilde{\mu}) = 1/\sqrt{2}\sum_{q_1q_2} p^{\lambda\mu}_{q_1q_2}(M)(u_{q_1}v_{q_2}-v_{q_1}u_{q_2})\left(T^{+}(q_1q_2\tilde{\mu})+T(q_1q_2\tilde{\mu})\right)+ \tag{33}$$

$$+\sum_{q_1q_2} p^{\lambda\mu}_{q_1q_2}(M)(u_{q_1}u_{q_2}+v_{q_1}v_{q_2})\,\mathcal{B}(q_1q_2\tilde{\mu}),$$

where

$$T^{+}(q_1q_2\tilde{\mu}) = -\frac{1}{\sqrt{1+\delta_{\mu,0}}}\sum_{\sigma_1\sigma_2}\delta_{\tilde{K}_1+\tilde{K}_2,\tilde{\mu}}\,\alpha^{+}_{\tilde{q}_1}\alpha^{+}_{\tilde{q}_2}\,\theta_{-\sigma_1-\sigma_2} \tag{34}$$

and

$$\mathcal{B}(q_1q_2\tilde{\mu}) = \sum_{\sigma_1\sigma_2}\delta_{\tilde{K}_1+\tilde{K}_2,\tilde{\mu}}\,\alpha^{+}_{\tilde{q}_1}\alpha_{\tilde{q}_2}\,\theta_{\sigma_1\sigma_2} \tag{35}$$

are the magnetic counterparts of electrical-type operators (15.1) and (15.2) and $p^{\lambda\mu}_{q_1q_2}(X)$ is the single-particle matrix element of $X\lambda$ transition. Since in our approach the long-range residual interaction is restricted to multipole forces only (and does not include the spin-multipole forces), the vibrations of doubly-even core are described by phonons (10) of electric type. As a result, the contribution of core polarization to the $\mathcal{M}'(X\lambda,\tilde{\mu})$ can be expressed in terms of phonon operators only for the $X=E$ case (see the second term in (31), where $L^{(E)}_{g}$ is the matrix element of $E\lambda$ transition between ground state and one-phonon state g in doubly-even core). The corresponding contribution for $X=M$ case is written in terms of two-quasiparticle operators (34).

Using (31) and (33) we can obtain the expressions for intrinsic matrix elements of transition operators:

$$\langle\Psi_{\nu'\gamma'_0}(\tilde{K}'_0)|\mathcal{M}'(X\lambda,\tilde{\mu})|\Psi_{\nu\gamma_0}(\tilde{K}_0)\rangle = \delta_{\tilde{K}'_0-\tilde{K}_0,\tilde{\mu}}\,\times$$

$$\times\sum_{\substack{sr\\s'r'}} C^{\nu'\gamma'_0}_{s'r'}C^{\nu\gamma_0}_{sr}\left(\tilde{p}^{\lambda\mu}_{ss'}(X)(u_su_{s'}\mp v_sv_{s'})\delta_{rr'}+\tilde{p}^{\lambda\mu}_{rr'}(X)(u_ru_{r'}\mp v_rv_{r'})\delta_{ss'}\right)k^{CC,+}_{\gamma_0\gamma'_0}$$

$$+\sum_{sri}\left(C^{\nu'\gamma'_0}_{sr}\mathcal{D}^{\nu\gamma_0}_{srg}k^{K_0}_{\mu}+\mathcal{D}^{\nu'\gamma'_0}_{srg}C^{\nu\gamma_0}_{sr}k^{K'_0}_{\mu}\right)L^{(X)}_{g}+ \tag{36}$$

$$+\sum_{\substack{sr\\s'r'g'\gamma\gamma'}}\mathcal{D}^{\nu\gamma_0}_{srg'}\mathcal{D}^{\nu'\gamma'_0}_{s'r'g'}\left(\tilde{p}^{\lambda\mu}_{ss'}(X)(u_su_{s'}\mp v_sv_{s'})\delta_{rr'}+\tilde{p}^{\lambda\mu}_{rr'}(X)(u_ru_{r'}\mp v_rv_{r'})\delta_{ss'}\right)k^{DD}_{\gamma\gamma'},$$

where, opposite to (32), the quantity $L_g^{(M)}$ is given by

$$L_g^{(M)} = \frac{\sqrt{1+\delta_{\mu,0}}}{2} \sum_{q_1 q_2} P_{q_1 q_2}^{\lambda\mu}(M) \left(u_{q_1} v_{q_2} - v_{q_1} u_{q_2} \right) \Psi_{q_1 q_2}^{g} \tag{37}$$

and in the right-hand side of (36) the upper and lower signs are valid for $X=E$ and $X=M$, respectively. Further,

$$k_{\gamma_0 \gamma_0'}^{CC} = \begin{cases} \frac{1}{\sqrt{2}} \left(\delta_{\sigma_{s_1},+1} \delta_{\sigma_{r_1},-1} - \gamma_0 \delta_{\sigma_{s_1},-1} \delta_{\sigma_{r_1},+1} \right), & if\ K_0=0,\ K_0' \neq 0 \\ \frac{1}{\sqrt{2}} \left(\delta_{\sigma_{s_1'},+1} \delta_{\sigma_{r_1'},-1} - \gamma_0' \delta_{\sigma_{s_1'},-1} \delta_{\sigma_{r_1'},+1} \right), & if\ K_0 \neq 0,\ K_0'=0 \\ 1 & ,\ in\ other\ cases \end{cases} \tag{38}$$

$$k_{\gamma\gamma'}^{DD} = \begin{cases} (\sqrt{2} k_\mu^{K_0'})^{-1} \left(\delta_{\sigma_{s_1},+1} \delta_{\sigma_{r_1},-1} - \gamma \delta_{\sigma_{s_1},-1} \delta_{\sigma_{r_1},+1} \right), & if\ K=0,\ K' \neq 0 \\ (\sqrt{2} k_\mu^{K_0})^{-1} \left(\delta_{\sigma_{s_1'},+1} \delta_{\sigma_{r_1'},-1} - \gamma' \delta_{\sigma_{s_1'},-1} \delta_{\sigma_{r_1'},+1} \right), & if\ K \neq 0,\ K=0 \\ k_\mu^{K_0'} k_\mu^{K_0} \left(1 + \delta_{K_0,0} \delta_{K_0',0} (1 - \delta_{\mu,0}) \right), & in\ other\ cases. \end{cases} \tag{39}$$

It should be noted that the shift operator of intrinsic angular momentum j^+ belongs to the magnetic type operator with $\lambda\mu$ =11 and, therefore, its matrix element presented in (A4), (A5) and (A6) can be determined by (36) for $X=M$.

7. Summary

The microscopic approach for description of low-lying states in odd-odd deformed nuclei is proposed. The approach is derived as a unification of the QPM [4] and the models where the rotational degrees of freedom as well as the Coriolis mixing and effective n-p interaction are included [6,29,30]. As a result, the approach takes into account the most important effects (n-p interaction between external nucleons and coupling with rotational and vibrational degrees of freedom of doubly even core) which are necessary for description of excitation energies as well as $E\lambda(M\lambda)$ -transitions in odd-odd deformed nuclei. It should be noted that the effects listed above are treated on the same microscopic footing. In particular, the n-p interaction responsible for the Gallagher-Moszkowski splitting and the Newby shift is shown to be a part of residual interaction which has been neglected earlier in the QPM.

Some comments on the approach have to be done. Firstly, the Pauli principle has to be allowed for. For this aim, the corresponding formulae for odd deformed nuclei [26] can be quite easy rewritten for the case of odd-odd nuclei. On the other hand, in refs. [35,36] it has been shown that the Pauli principle has to be taken into account simultaneously with coupling with multiphonon configurations since both these effects are of the same order and act opposite to each other. For the embedding of the latter effect would lead to tremendous complication of the approach, we did not allowed for both these effects and limited ourselves to the use of the simple procedure proposed in [26] for odd nuclei for indication of the most crucial violations of the Pauli principle. Secondly, two other modifications of the approach are very desirable: the exact exclusion of the "spurious" states caused by nonconservation of the particle number in an odd-odd nucleus (for an even-even core this problem is solved correctly [25]), and including the spin-multipole residual forces into consideration which may improve considerably the description of $M\lambda$ -transitions. Both the modifications are now in progress.

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Appendix A. Matrix elements of the Hamiltonian

In order to obtain the amplitudes $b_{\nu K}^{I\rho}$ in wave function (4) one must diagonalize the matrix of the Hamiltonian (1) in the basis of functions $|I^{\pi}MK\nu\rangle$ given by (5). The operator R_i in (5), representing the rotation by angle π around the second axis, changes the sign of K for the intrinsic state $\Psi_\nu(K^\pi)$. The special care has to be given to the $K=0$ case for which we have [32]:

$$R_i\,\Psi_{\nu\gamma}(K=0)=\gamma\,\Psi_{\nu\gamma}(K=0)\,, \tag{A1}$$

where $\gamma=\pm1$ and there is the condition

$$\gamma=(-1)^{I}. \tag{A2}$$

So, for $K=0$ the intrinsic function $\Psi_{\nu\gamma}(K=0)$ is the eigenvector of the operator R_i with eigenvalue γ. Therefore, in the paper, where it is needed, we ascribe the additional index γ to the intrinsic function. Following (A2) the rotational band based on the intrinsic state $\Psi_{\nu\gamma=+1}(K=0)$ can involve only the states with even values of I, while band based on $\Psi_{\nu\gamma=-1}(K=0)$ can have only odd values of I.

The matrix elements of total hamiltonian H in the basis of functions (5) are as follows

$$\langle I_1^{\pi_1}M_1K_1\nu_1|H|I_2^{\pi_2}M_2K_2\nu_2\rangle=\langle I_1^{\pi_1}M_1K_1\nu_1|H_{intr}+H_R+H_{cI}+H_j|I_2^{\pi_2}M_2K_2\nu_2\rangle, \tag{A3}$$

where

$$\langle I_1^{\pi_1}M_1K_1\nu_1|H_{intr}|I_2^{\pi_2}M_2K_2\nu_2\rangle=\delta_{I_1M_1K_1\nu_1,\,I_2M_2K_2\nu_2}\Big\{\eta_{\nu_1}(1-\delta_{K_1,0})+ \tag{A4}$$

$$+\tfrac{1}{2}\left[1+(-1)^{I}\gamma_1\right]\eta_{\nu_1\gamma_1}\cdot\delta_{\gamma_1,\gamma_2}\,\delta_{K_1,0}\Big\}$$

$$\langle I_1^{\pi_1}M_1K_1\nu_1|H_R|I_2^{\pi_2}M_2K_2\nu_2\rangle=\delta_{I_1M_1K_1\nu_1,\,I_2M_2K_2\nu_2}\cdot\frac{\hbar^2}{2\Phi}\left\{I_1(I_1+1)-K_1^2\right\}, \tag{A5}$$

$$\langle I_1^{\pi_1} M_1 K_1 \nu_1 | H_{CI} | I_2^{\pi_2} M_2 K_2 \nu_2 \rangle = -\,\delta_{I_1 M_1, I_2 M_2} \frac{\hbar^2}{2\Phi} \times$$

$$\times \Big\{ \delta_{K_1, K_2-1} \Big[\sqrt{(I_2+K_2)(I_2-K_2+1)} \langle \Psi_{\nu_2}(K_2) | j^+ | \Psi_{\nu_1}(K_1) \rangle +$$

$$+ (-1)^{I_2} \sqrt{\frac{I_2(I_2+1)}{2}}\, \delta_{K_1,0}\, \delta_{K_2,1} \langle \Psi_{\nu_2}(K_2) | j^+ | \Psi_{\nu_1 \gamma_1}(K_1=0) \rangle \Big] + \tag{A6}$$

$$+ \delta_{K_1, K_2+1} \Big[\sqrt{(I_2-K_2)(I_2+K_2+1)} \langle \Psi_{\nu_1}(K_1) | j^+ | \Psi_{\nu_2}(K_2) \rangle +$$

$$+ (-1)^{I_2} \sqrt{\frac{I_2(I_2+1)}{2}}\, \delta_{K_1,1}\, \delta_{K_2,0} \langle \Psi_{\nu_1}(K_1) | j^+ | \Psi_{\nu_2 \gamma_2}(K_2=0) \rangle \Big] \Big\},$$

$$\langle I_1^{\pi_1} M_1 K_1 \nu_1 | H_j | I_2^{\pi_2} M_2 K_2 \nu_2 \rangle = \delta_{I_1 M_1 K_1, I_2 M_2 K_2} \frac{\hbar^2}{2\Phi} \times$$

$$\times \sum_{\nu_i K_i} \Big[\langle \Psi_{\nu_1}(K_1) | j^+ | \Psi_{\nu_i}(K_i) \rangle \langle \Psi_{\nu_2}(K_2) | j^+ | \Psi_{\nu_i}(K_i) \rangle + \tag{A7}$$

$$+ \langle \Psi_{\nu_i}(K_i) | j^+ | \Psi_{\nu_1}(K_1) \rangle \langle \Psi_{\nu_i}(K_i) | j^+ | \Psi_{\nu_2}(K_2) \rangle \Big].$$

Appendix B. Notation for the intrinsic Hamiltonian

The functions used in Sec. 3 are derived within the QPM for the case when isoscalar and ispvector forces are taken into account:

$$X^{g}_{\tau} = (1+\delta_{\mu,0}) \sum_{q_1 q_2 \in \tau} \frac{f^{g}_{q_1 q_2} f^{\lambda\mu}_{q_1 q_2} u^{2}_{q_1 q_2} \varepsilon_{q_1 q_2}}{\varepsilon^{2}_{q_1 q_2} - \omega^{2}_{g}}, \tag{B1}$$

$$\mathcal{Y}^{g}_{\tau} = Y^{g}_{\tau} + Y^{g}_{-\tau} \left\{ \frac{1-(\varkappa_0^{(\lambda\mu)}+\varkappa_1^{(\lambda\mu)}) X^{g}_{\tau}}{(\varkappa_0^{(\lambda\mu)}-\varkappa_1^{(\lambda\mu)}) X^{g}_{-\tau}} \right\}^{2}, \tag{B2}$$

$$Y^{g}_{\tau} = (1+\delta_{\mu,0}) \sum_{q_1 q_2 \in \tau} \frac{f^{g}_{q_1 q_2} f^{\lambda\mu}_{q_1 q_2} u^{2}_{q_1 q_2} \varepsilon_{q_1 q_2} \omega_g}{(\varepsilon^{2}_{q_1 q_2} - \omega_g^{2})^{2}}, \tag{B3}$$

$$f^{g}_{q_1 q_2} = f^{\lambda\mu}_{q_1 q_2} - \delta_{q_1, q_2} \Gamma^{g\tau}_{q_1} / \gamma^{g}_{\tau}, \tag{B4}$$

where $\varepsilon_{q_1 q_2} = \varepsilon_{q_1} + \varepsilon_{q_2}$ is the energy of two-quasiparticle state $q_1 q_2$, ω_g is the phonon energy. Expressions for $\Gamma^{g\tau}_{q_1}$ and γ^{g}_{τ} can be found in [25,33].

Further,

$$\Gamma^{g}_{q_1 q_2} = \sqrt{\frac{2}{\mathcal{Y}^{g}_{\tau}}}\, f^{\lambda\mu}_{q_1 q_2} v_{q_1 q_2} \tag{B5}$$

and the amplitudes $\psi^{g}_{q_1 q_2}$ and $\varphi^{g}_{q_1 q_2}$ are normalized as

$$\sum_{q_1 q_2} (\psi^{g}_{q_1 q_2} \psi^{g'}_{q_1 q_2} - \varphi^{g}_{q_1 q_2} \varphi^{g'}_{q_1 q_2}) = 2\, \delta_{gg'}. \tag{B6}$$

In Sect. 4 the following notation has been used:

$$\tilde{\Gamma}^{g}_{q_1 q_2 \in \tau} = \sqrt{\frac{2}{\mathcal{Y}^{g}_{\tau}}}\, v_{q_1 q_2} \tilde{f}^{\lambda\mu}_{q_1 q_2} \left(1+\delta_{K,0}(1-\delta_{K_0,0})\right)^{-1/2} -$$

$$- \delta_{q_1, q_2} \delta_{\lambda\mu, 20} \sqrt{2}\cdot G_{\tau} u_{q_1} v_{q_1} \sum_{q \in \tau} (u^{2}_{q} - v^{2}_{q}) \sum_{j} (\psi^{20j}_{qq} + \varphi^{20j}_{qq}) \tag{B7}$$

where

$$\tilde{f}^{\lambda\mu}_{q_1q_2} = f^{\lambda\mu}_{q_1q_2} \begin{cases} 1 , & |K_{q_1}-K_{q_2}| = \mu \\ \sigma_{q_1} , & K_{q_1}+K_{q_2} = \mu , \end{cases} \tag{B8}$$

$$f^{\lambda\mu}_{q_1q_2} = \begin{cases} \langle q_1+|\hat{f}^{\lambda\mu}|q_2+\rangle = \langle q_1-|\hat{f}^{\lambda\mu}|q_2-\rangle , & |K_{q_1}-K_{q_2}| = \mu \\ \langle q_1+|\hat{f}^{\lambda\mu}|q_2-\rangle = \langle\langle q_1-|\hat{f}^{\lambda\mu}|q_2+\rangle , & K_{q_1}+K_{q_2} = \mu . \end{cases} \tag{B9}$$

The transition single-particle matrix elements $\tilde{p}^{\lambda\mu}_{qq'}(x)$ and $p^{\lambda\mu}_{qq'}(x)$ have the same form as (B8) and (B9), respectively, with the corresponding substitution of the operator of $X\lambda$ transition.

In the diagonal case $(r=r', s=s', \gamma=\gamma')$ the matrix element of n-p interaction $\langle rs|V_{np}|r's'\rangle_{\mu\gamma\gamma'}$ is written as

$$\begin{aligned} &\langle rs|V_{np}|rs\rangle_{\mu\gamma} = \\ &= \frac{1}{2}\Big\{\big(\langle r\pm s\pm|V_{np}|r\pm s\pm\rangle\tilde{u}_{rs} - \langle r\mp s\pm|V_{np}|r\mp s\pm\rangle\tilde{v}_{rs}\big)\times \\ &\quad\times\delta_{K_s+K_r,|K_0\pm\mu|} + \\ &+\big(\langle r\pm s\mp|V_{np}|r\pm s\mp\rangle\tilde{u}_{rs} - \langle r\mp s\mp|V_{np}|r\mp s\mp\rangle\hat{v}_{rs}\big)\times \\ &\quad\times\delta_{|K_s-K_r|,|K_0\pm\mu|} - \\ &-\gamma\big(\langle r\pm s\mp|V_{np}|r\mp s\pm\rangle\hat{u}_{rs} + \langle r\mp s\mp|V_{np}|r\pm s\pm\rangle\tilde{v}_{rs}\big)\times \\ &\quad\times\delta_{|K_s-K_r|,0}\Big\} , \end{aligned} \tag{B10}$$

where

$$\tilde{u}_{rs} = u_r^2u_s^2 + v_r^2v_s^2 , \quad \hat{v}_{rs} = u_r^2v_s^2 + v_r^2u_s^2 . \tag{B11}$$

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