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THE RPA AND THE RESTORATION OF TRANSLATION SYMMETRY OF ROTATING NUCLEUS HAMILTONIAN

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1. INTRODUCTION

The importance of the restoration of the total nucleus Hamiltonian symmetries spontaneously violated by the average nuclear field has been pointed out in literature several times (see, e.g., ref. [1,2]). But almost all up to date papers concerning this problem have worked with nonrotating nuclei and the restoration of translation symmetry of total Hamiltonian of rotating nucleus hasn't been discussed at all.

In refs.[3,4] and[5] the microscopic nuclear model was proposed describing the positive parity vibrational modes of nuclear motion near the yrast line within the random phase approximation (RPA) method based on the self Consistent Cranking Model (SCCM). In these refs. the attempts were made to clear up the relation between the positive parity solutions of RPA equations of motion and the spurious Goldstone modes connected with rotational symmetry of the total Hamiltonian of rotating nucleus. Negative parity vibrational states near the yrast line hasn't been investigated.

Taking into account the facts given above we have concluded that the translation symmetry violation in rotating Hamiltonian and its relation to the negative parity states near the yrast line deserve a separate investigation. Namely, this paper is devoted to the following problems:

i) to construct residual interactions consistently with the restoration of the translation symmetry of the rotating nucleus Hamiltonian of SCCM (sections 2,3);

ii) to determine structure and energies of the RPA phonons using the residual interactions obtained from restoration of the translation symmetry of nucleus Hamiltonian (sections 4, 5,6);

iii) to clear up the relation of the solutions of RPA equations of motion to the spurious states (Goldstone modes) connected with translation symmetry of the nuclear Hamiltonian (sections 4,5,6).

In solving the problems above we have followed the ideology of the SCCM+RPA method proposed by Marshalek[3] and Jaussen with Mikhailov[5] for a nucleus rotating around a fixed axis (stationary rotating). This method consists of two separate steps: In the first, the SCCM solution is found, which describes the yrast line states (using the Hartree-Fock-Bogolubov diaginalization (see, e.g.,[6])). In the next step, both the parity vibrations about the SCCM solutions are determined by the RPA. In this paper only the negative parity vibrations are treated (sections 4,5,6).

2. SYMMETRY OF THE SCCM HAMILTONIAN

In the SCCM the nucleus is supposed to rotate round the stable axis (axis x) and the number of nucleons has to be conserved. Therefore the Hamiltonian of the SCCM is taken in the form (see, e.g., [3,4]):

$$H' = \tilde{H} - \Omega J_x - \lambda \hat{N} \tag{2.1a}$$

$$\tilde{H} = e^{iJ_x\Omega t} H e^{-iJ_x\Omega t}, \tag{2.1b}$$

where H and $\tilde{H}$ are the total nucleus Hamiltonian in the lab. and intrinsic-fixed body system, respectively. J_x is an x-component of the total angular momentum, Ω is the rotating frequency, $\hat{N}$ is the particle number operator and the Lagrangian multiplies λ can be interpreted as the chemical potential. The total Hamiltonian H contains the deformed average nuclear field, pairing and the residual interactions. Solving the Hartree-Fock-Bogolubov equation, the SMMC Hamiltonian can be rewritten in the form [3,4]

$$H' = \langle\Omega|H|\Omega\rangle + \sum_j (E_j \alpha_j^+ \alpha_j + E_{\bar{j}} \alpha_{\bar{j}}^+ \alpha_{\bar{j}}) + \tag{2.2}$$
$$+ :H_{pair}: + V_{RES}(+) + V_{RES}(-),$$

where $\langle\Omega|H|\Omega\rangle$ denotes the mean value in quasiparticle vacuum (yrast state with given Ω), E_j are the quasiparticle energies, α_j^+ and α_j are the quasiparticle creation and annihilation operators, respectively. The symbol : : represents the normal product with respect to quasiparticle vacuum, H_{pair} is the pairing interactions. From the point of view of the space reflection symmetry, the residual interaction can be divided into two parts: a part containing the positive parity single particle operators($V_{RES}(+)$) and a part with the negative parity single particle ($V_{RES}(-)$).

The total nuclear Hamiltonian fulfils the following conservation laws:

$$[H, \hat{N}] = [H, \vec{J}] = [H, \vec{P}] = 0, \tag{2.3}$$

where $\vec{J}$ is the total spin operators, and $\vec{P}$ is the total linear momentum operator. Using the well-known commutation relations.

$$[J_k, J_\ell] = i\epsilon_{k\ell m} J_m$$
$$[J_k, \hat{N}] = [\tilde{P}_k, \hat{N}] = [\tilde{P}_k, \tilde{P}_\ell] = 0 \tag{2.4}$$
$$[J_k, \tilde{P}_\ell] = i\epsilon_{k\ell m} \tilde{P}_m$$
$$\tilde{P}_k = e^{iJ_x\Omega t} P_k e^{-iJ_x\Omega t}, \qquad k, \ell, m = x, y, z$$

from (2.3) it follows (see [3] and Appendix A):

$$[H, \hat{N}] = 0 \tag{2.5a}$$

$$[H', J_x] = 0, \quad [H', J_y] = -i\Omega J_z, \quad [H', J_z] = i\Omega J_y \tag{2.5b}$$

$$[H', \tilde{P}_x] = 0, \quad [H', \tilde{P}_z] = -i\Omega P_y, \quad [H', \tilde{P}_z] = i\Omega \tilde{P}_y. \tag{2.5c}$$

The conditions (2.5c) and (2.3) are used in the following sections to determine negative parity residual interactions.

3. RESTORATION OF THE TRANSLATION SYMMETRY

The violation of the translation invariance of nuclear Hamiltonian leads to the appearance of the admixtures connected with the center of mass motion in the odd parity states. Since the odd parity states near the yrast line are the subject of investigation in this paper, the nearest task is to construct the residual interactions which restore the translation invariance of the deformed average nuclear field. A possibility of the nonaxial average nuclear field is taken into account for the rotating nucleus.

According to [1] the separable residual interaction can be looked for in the form (we must have in mind that our nuclear average field rotates with the nucleus)

$$\tilde{H} = \tilde{H}_0 - \sum_{\mu=x,y,z} \frac{\kappa_\mu}{2} [\tilde{H}_0, \tilde{P}_\mu] [\tilde{H}_0, \tilde{P}_\mu] \tag{3.1a}$$

$$\tilde{H}_0 = e^{iJ_x\Omega t} H_0 e^{-iJ_x\Omega t},$$

where H_0 is the Hamiltonian of the average nuclear field with the pairing interaction (in lab.system), κ_μ are arbitrary constants for the time being, $\tilde{P}_\mu$ are the components of the nucleus total linear momentum (see sect.2). The requirement of the translation symmetry of nuclear Hamiltonian H can be expressed by the last commutator in (2.3). Substitution of (3.1) into (2.3) yields:

$$[\tilde{H}, \tilde{P}_\mu] = [\tilde{H}_0, P_\mu] - \sum_\nu \frac{\kappa_\nu}{2} \{[\tilde{H}_0, P_\nu], [[\tilde{H}_0, \tilde{P}_\nu], \tilde{P}_\mu]\} = 0, \qquad (3.2)$$

where the symbol { } denotes the anticommutator. From (3.2) one can see that if the double commutator is a c-number:

$$[[\tilde{H}_0, \tilde{P}_\nu], \tilde{P}_\mu] = c_\mu \delta_\mu \qquad (3.3)$$

it is possible to obtain the precise restoration of the nuclear Hamiltonian translation symmetry in the framework of the separable residual interaction, assuming $\frac{1}{\kappa_\nu} = c_\nu$ (see (3.2)). Unfortunately, the condition (3.3) holds only for the oscillator average nuclear field. Since the oscillator is rarely used as a nuclear average field, there is no possibility to restore precisely the translation symmetry for average fields usually used in nuclear theories. Because the problem of the nucleon motion in an atomic nucleus is never solved exactly, it is necessary for condition (2.3) to be fulfilled at the same stage of accuracy as the equation of motion is solved. In this paper the RPA is used for the description of the excited states near the yrast line, and therefore, the translation invariance condition (2.3) can be rewritten in the form:

$$[\tilde{H}, P_\mu] = [\tilde{H}, P_\mu]_{RPA} + \delta [\tilde{H}, P_\mu], \qquad (3.4)$$

where $[\tilde{H}, P_\mu]_{RPA}$ is the part of the commutator, which contains the boson operator in the linear order (see the next section and refs. /3,4/) and $\delta [\tilde{H}, P_\mu]_{RPA}$ is the remaining part of the commutator constructed from the second and higher order in boson powers. Neglecting the higher than linear terms in the boson expansion (3.4), one expressed the translation invariance condition in the RPA by:

$$[\tilde{H}, \tilde{P}_\mu]_{RPA} = 0 \qquad (\mu = x, y, z). \qquad (3.2')$$

Substituting (3.1) into (3.2') we obtain:

$$[\tilde{H}, \tilde{P}_\mu]_{RPA} = [\tilde{H}_0, \tilde{P}_\mu]_{RPA} - \sum_{\nu = x,y,z} \frac{\kappa_\nu}{2} \{[\tilde{H}_0, P_\nu], [\tilde{P}_\mu, [\tilde{H}_0, \tilde{P}_\nu]]\}_{RPA} =$$
$$= [\tilde{H}_0, \tilde{P}_\mu]_{RPA} - \sum_{\nu = x,y,z} \kappa_\nu \langle 0 | [\tilde{P}_\mu, [\tilde{H}_0, \tilde{P}_\nu]] | 0 \rangle \, [H_0, P_\nu]_{RPA} = 0, \qquad (3.5)$$

where $\langle 0|H|0\rangle$ denotes the boson (quasiparticle) vacuum (see /3/). From the reflection symmetry of the average nuclear field, one can get:

$$\langle 0 | [\tilde{P}_\mu, [\tilde{H}_0, P_\nu]] | 0 \rangle = c_\mu \delta_{\mu\nu}. \qquad (3.6)$$

Assuming $c_\mu = \frac{1}{\kappa_\mu}$, the RPA translation invariance condition (3.5) is automatically fulfilled.

Further, the structure of the nuclear average field will be discussed in detail. Suppose that the average field has the Saxon-Woods form /6/ with the following isotopic structure:

$$V = \sum_{i=1}^{Z} V^{(p)}(\vec{r}_i; \beta) + \sum_{i=1}^{A-Z} V^{(n)}(\vec{r}_i; \beta), \qquad (3.7)$$

where

$$V^{(n)}(\vec{r}; \beta) = V^{[0]}(\vec{r}; \beta) + V^{[1]}(\vec{r}; \beta)\tau_z + V^{(n)}_{\ell s}(\vec{r}; \beta), \qquad (3.8a)$$

$$V^{(p)}(\vec{r}; \beta) = V^{[0]}(\vec{r}; \beta) + V^{[1]}(\vec{r}; \beta)\tau_z + V^{(p)}_{\ell s}(\vec{r}; \beta) + V^{[c]}(\vec{r}; \beta). \qquad (3.8b)$$

In (3.7) and (3.8) symbol β denotes the set of deformation parameters, τ_z is the z-component of isospin for a given nucleon. Since the spinorbit interaction does not play the substantial role in the restoration of translation symmetry, it will be neglected in a further consideration. The restoration of the translation symmetry of the average nuclear field (3.8) can be made separately for the isoscalar, isovector and Coulomb interactions. Therefore, the residual interaction can be written in the form:

$$V = -\sum_{\substack{\nu = x,y,z \\ \tau = 0,1,c}} \frac{\kappa_\nu^{[\tau]}}{2} Q^+_\nu[\tau] \, Q_\nu[\tau], \qquad (3.9)$$

where

$$Q_\nu[0] \equiv [\sum_{i=1}^{A} V^{[0]}(\vec{r}_i; \beta), \tilde{P}_\nu], \qquad (3.10)$$

$$Q_\nu[1] \equiv [\sum_{i=1}^{A} V^{[1]}(r_i; \beta)\tau_z(i), \tilde{P}_\nu],$$

$$Q_\nu[c] \equiv [\sum_{i=1}^{Z} V^{[c]}(r_i; \beta), \tilde{P}_\nu].$$

In accordance with (3.6) the strength constants $\kappa_\nu^{[\tau]}$ are

$$\frac{1}{\kappa_\nu^{[\tau]}} = \langle 0 | [\tilde{P}_\nu, Q_\nu^{[\tau]}] | 0 \rangle. \qquad (3.11)$$

The residúal interactions given by (3.9), (3.10), (3.11) are further used for constructing the negative parity vibrational states near the yrast line.

4. RPA HAMILTONIAN AND THE CONSERVATION LAWS

Introducing the two-quasiparticle boson creation and annihilation operators as in papers [3,4]: $b^+_{k\ell} = a^+_k a^+_\ell$, $b^+_{k\bar\ell} = i a^+_k a^+_{\bar\ell}$,

$b^+_{\bar k\bar\ell} = i\, a^+_{\bar k} a^+_{\bar\ell}$ and taking into account the behaviour of the single-particle operators (involved in the residual interactions) under the space reflections and the rotation around the x-axis by angle π, the nuclear Hamiltonian of the SCCM (2.2), can be divided in the framework of the RPA (i.e., up to the second order in the boson expansion) into mutually commuting parts (see [3,7,8])

$$H' = \langle\Omega|H|\Omega\rangle + H\binom{\gamma_x=+1}{\gamma_P=+1} + H\binom{\gamma_x=-1}{\gamma_P=+1} + H\binom{\gamma_x=+1}{\gamma_P=-1} + H\binom{\gamma_x=-1}{\gamma_P=-1}, \quad (4.1)$$

where γ_x and γ_P are quantum numbers characterizing the properties of the single-particle operators, involved in residual interactions, under a rotation of angle π around the x-axis and space reflection, respectively. Because the problem of searching of eigenvalues and corresponding eigenvectors of the Hamiltonian (4.1) for the positive parity part ($\gamma_P = +1$) has been solved in refs. [3,4,5], only the negative parity part ($\gamma_P = -1$) of (4.1) is treated in this paper. With respect to (3.9) the negative parity part of the Hamiltonian (4.1) can be expressed in the form:

$$H\binom{\gamma_x=+1}{\gamma_P=-1} = \sum_{ik} E_{i\bar k}\, b^+_{i\bar k} b_{i\bar k} - \sum_{\tau=0,1,c} \frac{\kappa_x^{[\tau]}}{2} Q_x^{(1)^+}[\tau]\, Q_x^{(1)}[\tau], \quad (4.2a)$$

$$H\binom{\gamma_x=-1}{\gamma_P=-1} = \frac{1}{2}\sum_{ik}\left(E_{ik} b^+_{ik} b_{ik} + E_{\bar i\bar k} b^+_{\bar i\bar k} b_{\bar i\bar k}\right) - \sum_{\substack{\tau=0,1,c\\ \mu=y,z}} \frac{\kappa_\mu^{[\tau]}}{2} Q_\mu^{(1)^+}[\tau]\, Q_\mu^{(1)}[\tau], \quad (4.2b)$$

where $E_{ik} = E_i + E_k$, $Q^{(1)}_\mu[\tau]$ ($\tau = 0, 1, c$; $\mu = x, y, z$) are the linear boson parts of the single-particle operators (3.10). The explicit forms of $Q^{(1)}_\mu[\tau]$, together with the linear boson part of linear momentum components and their conjugated coordinates are given in Appendix B.

In the Hartree-Fock-Bogolubov representation the Hamiltonian H' does not contain the linear boson part, therefore up to the second order in bosons (see [3,4])

$$H' = \langle 0|H|0\rangle + H'^{(2)}. \quad (4.3)$$

Substituting (4.3), (4.2) into (2.4) and (2.5c) and using the boson expansions for the angular momentum components given in ref. [3] one can get:

$$[H'^{(2)}, P_x^{(1)}] = 0,\quad [H'^{(2)}, P_y^{(1)}] = -i\Omega P_z^{(1)},\quad [H'^{(2)}, P_z^{(1)}] = i\Omega P_y^{(1)}, \quad (4.4a)$$

$$[P_i^{(1)}, P_j^{(1)}] = [N^{(1)}, P_j^{(1)}] = [J_i^{(1)}, P_j^{(1)}] = [J_x^{(1)}, J_j^{(1)}] = 0, \quad (4.4b)$$

$$[J_y^{(1)}, J_z^{(1)}] = i\langle 0|J_x|0\rangle. \quad (4.4c)$$

Since all parts of the Hamiltonian H' (4.1) mutually commute, the problem of searching its eigenvalues and corresponding eigenvectors can be solved independently for each part of the RPA Hamiltonian (see [3,7])

$$[H', \mathcal{P}_\lambda] = i\omega_\lambda^2 \mathcal{X}_\lambda, \quad (4.5a)$$

$$[H', \mathcal{X}_\lambda] = -i\mathcal{P}_\lambda, \quad (4.5b)$$

$$[\mathcal{X}_\lambda, \mathcal{P}_{\lambda'}] = i\delta_{\lambda\lambda'}, \quad (4.5c)$$

where $\mathcal{X}_\lambda$ and $\mathcal{P}_\lambda$ represent canonical conjugate coordinates and momenta in the state λ with energy ω_λ ($\mathcal{X}_\lambda$ and $\mathcal{P}_\lambda$ are supposed in the linear boson order). In this canonical representation the Hamiltonian (4.1) can be rewritten (see [7]):

$$H' = \frac{1}{2}\sum_{\substack{\lambda\\(\omega_\lambda\neq 0)}} \left(\mathcal{P}_\lambda^2 + \omega_\lambda^2 \mathcal{X}_\lambda^2\right) + \frac{1}{2}\sum_{\substack{\lambda_0\\(\omega_{\lambda_0}=0)}} \mathcal{P}_{\lambda_0}^2. \quad (4.6)$$

From the comparison of (4.4a) with the RPA equation of motion (4.5) it follows

$$\mathcal{P}_{\lambda_0} = \sqrt{g_x}\, P_x^{(1)}, \quad (4.7)$$

where the constant g_x ("mass" parameter) will be discussed later (see section 5). Therefore, the negative parity part of Hamiltonian (4.1) has the form:

$$H'_{negat.} = \frac{1}{2}\sum_{\substack{\lambda\\ \omega_\lambda\neq 0}} \left(\mathcal{P}_\lambda^2 + \omega_\lambda^2 \mathcal{X}_\lambda^2\right) + \frac{1}{2} g_x P_x^{(1)^2} + H_{c.m.}(P_y^{(1)}, P_z^{(1)}, X_y^{(1)}, X_z^{(1)}), \quad (4.8)$$

where $H_{c.m.}(P_y^{(1)}, P_z^{(1)}, X_y^{(1)}, X_z^{(1)})$ is the part of Hamiltonian connected with the center of mass motion modes $(X_y^{(1)}, P_y^{(1)})$, $(X_z^{(1)}, P_z^{(1)})$ which cannot be interpreted as the solutions of RPA equations of motion. These modes are discussed in detail in section 5.2.

The structure of the odd parity vibrational states (i.e., canonical conjugated coordinates $\mathcal{X}_\lambda$ and momenta $\mathcal{P}_\lambda$ and corresponding energies ω_λ) near the yrast line for a given angular momentum is determined in the following sections.

5. RPA EQUATIONS FOR THE NAGETIVE PARITY PART OF THE HAMILTONIAN H

The RPA equations (4.5) with negative parity part of the Hamiltonian H' can be solved by the method proposed by Kvasil et al.[7], where the way of extraction of the spurious states from the solution is also described.

5.1. Diagonalization of $H_{(+)} \equiv H(\gamma_x = +1, \gamma_p = -1)$

The linear boson part of canonical conjugated coordinates and momenta can be looked for in the form

$$\mathcal{X}_\lambda = -i \sum_{ik} \mathcal{X}_{ik}^{(\lambda)} (b_{i\bar{k}}^{+} - b_{i\bar{k}}), \qquad (5.1a)$$

$$\mathcal{P}_\lambda = \sum_{ik} \mathcal{P}_{ik}^{(\lambda)} (b_{i\bar{k}}^{+} + b_{i\bar{k}}), \qquad (5.1b)$$

$$[\mathcal{X}_\lambda, \mathcal{P}_{\lambda'}] = 2i \sum_{ik} \mathcal{X}_{ik}^{(\lambda)} \mathcal{P}_{ik}^{(\lambda')} = i\delta_{\lambda\lambda'}. \qquad (5.1c)$$

Substituting expressions (5.1a), (5.1b), and (4.2a) into (4.5) and using the commutation relation of boson operators (see [3,4]), one can obtain equations of motion in the matrix representation:

$$\mathcal{X}_{ik}^{(\lambda)} = 2 \sum_{r=0,1,c} \kappa_x^{[r]} A_\lambda^{[r]} \frac{\tilde{q}_{ik}^{x}[r]}{E_{i\bar{k}}^2 - \omega_\lambda^2}, \qquad (5.2a)$$

$$\mathcal{P}_{ik}^{(\lambda)} = 2 \sum_{r=0,1,c} \kappa_x^{[r]} A_\lambda^{[r]} \frac{E_{i\bar{k}} \tilde{q}_{ik}^{x}[r]}{E_{i\bar{k}}^2 - \omega_\lambda^2}, \qquad (5.2b)$$

where

$$A_\lambda = \sum_{\ell p} \mathcal{P}_{\ell p}^{(\lambda)} \tilde{q}_{\ell p}^{x}[r]. \qquad (5.3)$$

The summation in (5.2) is over the isoscalar, isovector, and Coulomb part (see (4.2a) and Appendix B). Constructing the corresponding linear combinations of equations (5.2a,b) and using the way described in ref.[7] one can get the following system of the linear homogeneous algebraic equations for three unknowns $A_\lambda^{[r]}$ $(r = 0,1,c)$

$$\omega_\lambda^2 \sum_{r=0,1,c} \kappa_x^{[r]} A_\lambda^{[r]} \Big(\sum_{r'=0,1,c} W^{(\lambda)}_{q^x[r]\, q^x[r']}\Big) = 0,$$

$$\sum_{r=0,1,c} \kappa_x^{[r]} A_\lambda^{[r]} \Big(S^{(\lambda)}_{q^x[c]\, q^x[r]}(+) - \frac{\delta_{rc}}{2\kappa_x^{[c]}}\Big) = 0, \qquad (5.4)$$

$$\sum_{r=0,1,c} \kappa_x^{[r]} A_\lambda^{[r]} \Big(S^{(\lambda)}_{q^x[c]\, q^x[r]}(+) - \frac{\delta_{rc}}{2\kappa_x^{[c]}}\Big) = 0,$$

where

$$S_{a,b}^{(\lambda)}(+) \equiv \sum_{ik} \frac{E_{i\bar{k}}\, \tilde{a}_{ik}\, \tilde{b}_{ik}}{E_{i\bar{k}}^2 - \omega_\lambda^2}, \qquad W_{a,b}^{(\lambda)} = \sum_{ik} \frac{\tilde{a}_{ik}\, \tilde{b}_{ik}}{E_{i\bar{k}}(E_{i\bar{k}}^2 - \omega_\lambda^2)}, \qquad (5.5)$$

The conditions for existence of nontrivial solutions of the equation system (5.4) lead to the secular equations for the eigenvalues ω_λ in the form

$$\omega_\lambda^2 |D^{(+)}(\omega_\lambda)| = 0, \qquad (5.6)$$

where $|D^{(+)}(\omega_\lambda)|$ is the determinant of the equation system (5.4) for nonspurious solutions. The equation system (5.4) has to be solved taking into account the condition (5.1c).

The strength constants $\kappa_x^{[r]}$ $(r = 0,1,c)$ are determined by eq. (3.11) which can be rewritten by the quasiparticle matrix elements $\tilde{q}_{ik}^{x}[r]$ using the condition (4.4a): $[H'^{(2)}, P_x^{(1)}] = 0$ in the form:

$$\frac{1}{\kappa_x^{[r]}} = \sum_{r'=0,1,c} \Big(\sum_{ik} \frac{\tilde{q}_{ik}^{x}[r]\, \tilde{q}_{ik}^{x}[r']}{E_{i\bar{k}}}\Big). \qquad (5.7)$$

As the Hamiltonian $H_{(+)}$ commutes with the x-component of the total linear momentum $P_x^{(1)}$ (see eq. (4.4a)), among the solutions of the equation system (5.4) the spurious one with $\omega_\lambda = 0$ appears, which is connected with the center of mass motion along the x-axis (see (4.7)). The "mass" parameter g_x connected with this spurious mode can be determined from the normalization condition (5.1c). Using (4.7), (5.2), (5.3), and (5,7), one gets for mode:

$$\mathcal{X}_{ik}(\omega_{\lambda_0}=0)=\sqrt{g_x}\,\frac{\tilde{p}^x_{ik}}{E_{i\bar{k}}};\quad \mathcal{P}_{ik}(\omega_{\lambda_0}=0)=\sqrt{g_x}\,\tilde{p}^x_{ik}. \tag{5.8}$$

Then from (5.1c) the "mass" parameter can be obtained

$$\frac{1}{g_x}=2\sum_{ik}\frac{\tilde{p}^x_{ik}\,\tilde{p}^x_{ik}}{E_{i\bar{k}}}. \tag{5.9}$$

The completeness of all solutions $(\mathcal{X}_\lambda,\mathcal{P}_\lambda)$ of the $H_{(+)}$ Hamiltonian can be expressed by

$$\begin{aligned} b^+_{\ell\bar{p}} &= i\sum_\lambda (b^+_{\ell\bar{p}},\mathcal{X}_\lambda]\,\mathcal{P}_\lambda+[\mathcal{P}_\lambda,b^+_{\ell\bar{p}}]\,\mathcal{X}_\lambda)+\\ &+i([b^+_{\ell\bar{p}},X^{(1)}_x]\,P^{(1)}_x+[P^{(1)}_x,b^+_{\ell\bar{p}}]\,X^{(1)}_x). \end{aligned} \tag{5.10}$$

5.2. Diagonalization of $H_{(-)}\equiv H\begin{pmatrix}\gamma_x=-1\\ \gamma_p=-1\end{pmatrix}$

The canonical conjugated coordinates and momenta for this part of the Hamiltonian (4.1) are assumed to be in the form (linear boson terms only):

$$\mathcal{X}_\lambda=\sum_{ik}(\mathcal{X}^{(\lambda)}_{ik}(b^+_{ik}+b_{ik})+\tilde{\mathcal{X}}^{(\lambda)}_{ik}(b^+_{i\bar{k}}+b_{i\bar{k}})), \tag{5.11a}$$

$$\mathcal{P}_\lambda=i\sum_{ik}(\mathcal{P}^{(\lambda)}_{ik}(b^+_{ik}-b_{ik})+\tilde{\mathcal{P}}^{(\lambda)}_{ik}(b^+_{i\bar{k}}-b_{i\bar{k}})), \tag{5.11b}$$

$$[\mathcal{X}_\lambda,\mathcal{P}_\lambda]=4i\sum_{ik}(\mathcal{X}^\lambda_{ik}\mathcal{P}^\lambda_{ik}+\tilde{\mathcal{X}}^\lambda_{ik}\tilde{\mathcal{P}}^\lambda_{ik}). \tag{5.11c}$$

Similarly as in the preceding section the RPA equations of motion (4.5) for the part $H_{(-)}$ of the Hamiltonian (4.1) can be expressed in the matrix representation:

$$\mathcal{X}^{(\lambda)}_{ik}=\sum_{r=0,1,c}\kappa^{[r]}_y B^{[r]}_\lambda\frac{E_{ik}\,\tilde{q}^y_{ik}[r]}{E^2_{ik}-\omega^2_\lambda}+\sum_{r=0,1,c}\kappa^{[r]}_z C^{[r]}_\lambda\frac{\tilde{q}^z_{ik}[r]}{E^2_{ik}-\omega^2_\lambda}, \tag{5.12a}$$

$$\mathcal{P}^{(\lambda)}_{ik}=\sum_{r=0,1,c}\kappa^{[r]}_y B^{[r]}_\lambda\frac{\omega^2_\lambda\,\tilde{q}^y_{ik}[r]}{E^2_{ik}-\omega^2_\lambda}+\sum_{r=0,1,c}\kappa^{[r]}_z C^{[r]}_\lambda\frac{E_{ik}\tilde{q}^z_{ik}[r]}{E^2_{ik}-\omega^2_\lambda}, \tag{5.12b}$$

where the unknowns $B^{[r]}_\lambda$, $C^{[r]}_\lambda$ $(r=0,1,c)$ are introduced

$$\begin{aligned} B^{[r]}_\lambda &= \sum_{\ell p}(\tilde{q}^y_{\ell p}[r]\,\mathcal{X}^{(\lambda)}_{\ell p}-\tilde{\tilde{q}}^y_{\ell p}[r]\,\tilde{\mathcal{X}}^{(\lambda)}_{\ell p}),\\ C^{[r]}_\lambda &= \sum_{\ell p}(\tilde{q}^z_{\ell p}[r]\,\mathcal{P}^{(\lambda)}_{\ell p}+\tilde{\tilde{q}}^z_{\ell p}[r]\,\mathcal{P}^{(\lambda)}_{\ell p}). \end{aligned} \tag{5.13}$$

For $\tilde{\mathcal{X}}^{(\lambda)}_{ik}$ and $\tilde{\mathcal{P}}^{(\lambda)}_{ik}$ a relations similar to (5.12) can be obtained. Substitution of (5.12) into (5.13) leads to the homogeneous system of linear equations for six unknowns $B^{[r]}_\lambda$ and $C^{[r]}_\lambda(r=0,1,c)$:

$$\sum_{r'=0,1,c}\kappa^{[r]}_y B^{[r']}_\lambda(S^{(\lambda)}_{q^y[r]q^y[r']}(-)-\frac{\delta_{rr'}}{\kappa^{[r]}_y})+\sum_{r'=0,1,c}\kappa^{[r']}_z C^{[r']}_\lambda U^{(\lambda)}_{q^y[r]q^z[r']}=0, \tag{5.14}$$

$$\omega^2_\lambda\sum_{r'=0,1,c}\kappa^{[r']}_y B^{[r']}_\lambda U^{(\lambda)}_{q^z[r]q^y[r']}+\sum_{r'=0,1,c}\kappa^{[r']}_z C^{[r']}_\lambda(S^{(\lambda)}_{q^z[r]q^z[r']}(-)-\frac{\delta_{rr'}}{\kappa^{[r]}_z})=0$$

with the corresponding secular equation for the eigenvalues ω_λ. In (5.14) the following assignment is made

$$\begin{aligned} S^{(\lambda)}_{ab}(-) &= S_{ab}(\omega_\lambda)=\sum_{ik}(\frac{E_{ik}\tilde{a}_{ik}\tilde{b}_{ik}}{E^2_{ik}-\omega^2_\lambda}\pm\frac{E_{i\bar{k}}\tilde{\tilde{a}}_{ik}\tilde{\tilde{b}}_{ik}}{E^2_{i\bar{k}}-\omega^2_\lambda}),\\ U^{(\lambda)}_{ab} &= U_{ab}(\omega_\lambda)=\sum_{ik}(\frac{\tilde{a}_{ik}\tilde{b}_{ik}}{E^2_{ik}-\omega^2_\lambda}\pm\frac{\tilde{\tilde{a}}_{ik}\tilde{b}_{ik}}{E^2_{i\bar{k}}-\omega^2_\lambda}). \end{aligned} \tag{5.15}$$

The sign (-) holds in the case when one of the quantities $\tilde{a}$ or $\tilde{b}$ is $\tilde{q}^y[r]$ for all r.

If it is assumed that all conservation laws are exhausted by (2.5), $H_{(-)}$ (for $\Omega\neq 0$) has no Goldstone modes connected with zero energy $\omega_\lambda=0$.

The strength constants $\kappa^{[r]}_y$ and $\kappa^{[r]}_z$ in (5.14) are given by (3.11). Using the symmetry conditions (4.4a) these constants can be expressed by means of the quasiparticle matrix elements:

$$\frac{1}{\kappa_y^{[r]}} = \sum_{r'=0,1,c} S_{q^y[r],q^y[r']}(\omega = \Omega) = \Omega \sum_{r'=0,1,c} U_{q^y[r]\,q^z[r']}(\omega = \Omega), \qquad (5.16)$$

$$\frac{1}{\kappa_z^{[r]}} = \sum_{r'=0,1,c} S_{q^z[r]q^z[r']}(\omega = \Omega) = \Omega \sum_{r'=0,1,c} U_{q^z[r]\,q^y[r']}(\omega = \Omega).$$

The only question which remains uncleared for $H_{(-)}$ is wheather the canonical conjugated coordinates and linear momenta: $(X_y^{(1)}, P_y^{(1)})$, $(X_z^{(1)}, P_z^{(1)})$ of the nucleus center of mass motion in the rotation plane represent the spurious states of the Hamiltonian $H_{(-)}$. From comparison of (4.4a) with (4.5) one could expect that $(P_y^{(1)}, P_z^{(1)})$ is a mode of $H_{(-)}$ with energy $\omega = \pm\Omega$ (similarly as the y - and z -components of angular momentum were treated in paper of Marshalek[3]). But since $[P_y^{(1)}, P_z^{(1)}] = 0$, such a mode would be non-normalisable. In can be shown that for the same reason neither $(X_y^{(1)}, P_y^{(1)})$ nor $(X_z^{(1)}, P_z^{(1)})$ represents the RPA mode of Hamiltonian $H_{(-)}$ for $\Omega \neq 0$. In the case of $\Omega = 0$ (non-rotational nucleus) all linear momentum components commute with the Hamiltonian H' (see (4.4a)), and therefore $(X_y^{(1)}, P_y^{(1)})$ and $(X_z^{(1)}, P_z^{(1)})$ are zero spurious modes of $H_{(-)}$. Although the $(X_y^{(1)}, P_y^{(1)})$ and $(X_z^{(1)}, P_z^{(1)})$ don't belong to the ansamble of all RPA solutions with Hamiltonian $H_{(-)}$ in general case $\Omega \neq 0$, we give the proof of the orthogonality of the modes $(X_y^{(1)}, P_y^{(1)})$ and $(X_z^{(1)}, P_z^{(1)})$ to all solutions of the equation system (5.14).

Using the Jacobi identity one can obtain

$$[\mathcal{X}_\lambda, P_z^{(1)}] = (\text{see (4.5a)}) = \frac{1}{i\omega_\lambda^2}[[H_{(-)}, \mathcal{P}_\lambda], P_z^{(1)}] =$$
$$= -\frac{1}{i\omega_\lambda^2}[[P_z^{(1)}, H_{(-)}], \mathcal{P}_\lambda] - \frac{1}{i\omega_\lambda^2}[[\mathcal{P}_\lambda, P_z^{(1)}], H_{(-)}] =$$
$$= -\frac{1}{i\omega_\lambda^2}[[P_z^{(1)}, H_{(-)}], \mathcal{P}_\lambda] = (\text{see (4.4a)}) = \qquad (5.17a)$$
$$= -\frac{\Omega}{\omega_\lambda^2}[P_y^{(1)}, \mathcal{P}_\lambda] = (\text{see (4.5b)}) = -\frac{\Omega}{i\omega_\lambda^2}[P_y^{(1)}, [H', \mathcal{X}_\lambda]] =$$
$$= \frac{\Omega}{i\omega_\lambda^2}[\mathcal{X}_\lambda)[P_y^{(1)}, H_{(-)}]] + \frac{\Omega}{i\omega_\lambda^2}[H_{(-)}, [\mathcal{X}_\lambda, P_y^{(1)}]] =$$
$$= \frac{\Omega}{i\omega_\lambda^2}[\mathcal{X}_\lambda, [P_y^{(1)}, H_{(-)}]] = (\text{see (4.4a)}) = \frac{\Omega^2}{\omega_\lambda^2}[\mathcal{X}_\lambda, P_z^{(1)}].$$

In the same way we can obtain:

$$(1 - \frac{\Omega^2}{\omega_\lambda^2})[\mathcal{P}_\lambda, P_y^{(1)}] = 0, \qquad (5.17b)$$

$$(1 - \frac{\Omega^2}{\omega_\lambda^2})[\mathcal{X}_\lambda, X_y^{(1)}] = 0, \qquad (5.17c)$$

$$(1 - \frac{\Omega^2}{\omega_\lambda^2})[\mathcal{P}_\lambda, X_z^{(1)}] = 0. \qquad (5.17d)$$

Since $H_{(-)}$ has no solution with $\omega_\lambda = \pm\Omega$, in general case from (5.17), the orthogonality of all RPA solutions of $H_{(-)}$ to the $(X_y^{(1)}, P_y^{(1)})$ and $(X_z^{(1)}, P_z^{(1)})$ follows. We see, that the total space of canonical conjugated coordinates and momenta with $\gamma_x = -1$ and $\gamma_p = -1$ is a sum of the space of the normal modes (RPA solutions) $(\mathcal{X}_\lambda, \mathcal{P}_\lambda)$ of $H_{(-)}$ and the four-dimensional space of $(X_y^{(1)}, P_y^{(1)})$, $(X_z^{(1)}, P_z^{(1)})$. Therefore the completeness of this total space can be expressed by:

$$b_{\ell p}^{+} = i \sum_\lambda ([\mathcal{P}_\lambda, b_{\ell p}^{+}]\mathcal{X}_\lambda + [b_{\ell p}^{+}, \mathcal{X}_\lambda]\mathcal{P}_\lambda) +$$
$$+ i([P\ \ , b\ \ ]X\ \ + [b\ \ , X\ \ ]P\ \ ) \qquad (5.18)$$
$$b_{\overline{\ell p}}^{+} = i \sum ([\mathcal{P}_\lambda, b_{\overline{\ell p}}^{+}]\mathcal{X}_\lambda + [b_{\overline{\ell p}}^{+}, \mathcal{X}_\lambda]\mathcal{P}_\lambda) +$$
$$+ i([P_y^{(1)}, b_{\ell p}^{+}]X_y^{(1)} + [b_{\ell p}^{+}, X_y^{(1)}]P_y^{(1)}) + i([P_z^{(1)}, b_{\ell p}^{+}]X_z^{(1)} + [X_z^{(1)}, b_{\ell p}^{+}]P_z^{(1)}),$$

5.3. Parameters of the SCCM+RPA Method

Our SCCM+RPA treatment of the vibrational states near the yrast line has no free parameters (we don't speak about parameters of the average nuclear field). The form and strength constants of the residual interactions are wholly determined by the requirement of the restoration of the spontaneously broken symmetries of the average field. As is shown in the preceding sections, the odd part of the residual interaction is unambiguously given by the translation symmetry of the total nuclear Hamiltonian. The corresponding strength constants can be determined in terms of relations (5.7), (5.19).

6. THE EIGENVECTORS OF THE NUCLEUS HAMILTONIAN

In section 5 the procedure for searching the negative parity part ($\gamma_p=-1$) of the Hamiltonian (4.1) eigenvalues and the corresponding pair of conjugated coordinates $\mathcal{X}_\lambda$ and momenta $\mathcal{P}_\lambda$ is given. For the states with $\omega_\lambda \neq 0$ it is possible to introduce the phonon creation and annihilation operators /7/

$$\mathcal{C}^+_\lambda = \frac{1}{\sqrt{2}}\left(\sqrt{\omega_\lambda}\,\mathcal{X}_\lambda - \frac{i}{\sqrt{\omega_\lambda}}\mathcal{P}_\lambda\right), \tag{6.1}$$

$$[\mathcal{C}_\lambda, \mathcal{C}^+_{\lambda'}] = \delta_{\lambda\lambda'}.$$

Taking into account the fact that the modes $(X^{(1)}_y, P^{(1)}_y)$ and $(X^{(1)}_z, P^{(1)}_z)$ do not belong to the normal modes $(\mathcal{X}_\lambda, \mathcal{P}_\lambda)$ and that the conditions (4.4a) must be fulfilled, one can rewrite the negative parity part of (4.1) in the form:

$$H'_{\mathbf{negat.}} = \sum_{\substack{\lambda \\ (\omega_\lambda \neq 0)}} \omega_\lambda \left(\mathcal{C}^+_\lambda \mathcal{C}_\lambda + \frac{1}{2}\right) + \frac{1}{2} g_x P^{(1)^2}_x - \Omega\left(X^{(1)}_y P^{(1)}_z - X^{(1)}_z P^{(1)}_y\right), \tag{6.2}$$

where we have explicitly extracted the terms connected with the center of mass motion.

The positive parity ($\gamma_p=+1$) part of the Hamiltonian is expressed in a similar way (see /3,4/) in terms of the phonon operators for positive parity.

Since the Hamiltonian H' commutes with J_x-component of the angular momentum, with operator of the number of nucleons (for protons and neutrons separately), and with P_x-component of the linear momentum (see (4.4a)), the wave function describing the nucleus state near the yrast line is characterized by quantum numbers $\{n_{\lambda(-)}\}$, $\{n_{\lambda(+)}\}$, N, Z, P_x, J_x, where $n_{\lambda(+)}$ and $n_{\lambda(-)}$ are numbers of both parity phonons of positive and negative signature, respectively. For the states near the yrast line with high spin it is possible to add to these quantum numbers the total angular momentum J. Using the method proposed in paper of Marshalek /3/ for constructing the eigenstates of H', which generalizes the RPA results for collective modes of H', one can write:

$$|\{n_{\lambda(+)}\}\{n_{\lambda(-)}\}, N, Z, P_x, I, M\rangle = \prod_{\lambda(+)} \frac{(\mathcal{C}^+_\lambda)^{n_{\lambda(+)}}}{\sqrt{n_{\lambda(+)}!}} \prod_{\lambda(-)} \frac{(\mathcal{C}^+_\lambda)^{n_{\lambda(-)}}}{\sqrt{n_{\lambda(-)}!}} \times$$

$$\times \frac{e^{i(N-N_0)\chi_n}}{\sqrt{2\pi}} \cdot \frac{e^{i(Z-Z_0)\chi_p}}{\sqrt{2\pi}} \cdot \frac{e^{iP_x X}}{\sqrt{2\pi}} \cdot \frac{e^{i(J-J_0)\phi}}{\sqrt{2\pi}} \times \tag{6.3}$$

$$\times \frac{(\Gamma^+)^{J-M}}{\sqrt{(J-M)!}} |0, 0, N_0, Z_0, 0, J_0, J_0\rangle,$$

where χ_n, χ_p, ϕ and X_x are the variables conjugated to the $\hat{N}$, $\hat{Z}$, $\hat{J}$ and $\hat{P}_x$, respectively ($\hat{N}$ is the neutron number operator, and $\hat{Z}$ - the proton number operator). Notation in (6.3) is the same as in ref. /3/ The ket $|0, P, N_0, Z_0, 0, J_0, J_0\rangle$ describes the yrast line state with the angular momentum J_0 (with projection J_0 onto x-axis), without the translation motion of nucleus along x-axis and for nucleus with N_0 neutrons and Z_0 protons. In analogy with /3/ we assume the wave function $|\{n_{\lambda(-)}\}, \{n_{\lambda(+)}\}, N, Z, P_x, J, M\rangle$ to be invariant under rotation of angle π about the x-axis. This assumption leads to the condition (see /3/):

$$(-1)^{J + \sum_{\lambda(-)} n_{\lambda(-)}} = 1. \tag{6.4}$$

This allows us to conclude: if a nucleus state contains the even number of the negative signature phonons, the total angular momentum must be even, and if the state is formed by the odd number of negative signature phonons, the total momentum J is odd. This conclusion is independent of parity of phonons.

7. SUMMARY

The way of determination of the negative parity part of the residual interaction given in section 3 of this paper is general from the point of view of the form of nuclear average field. It can be used for the axial symmetry form as well as for the non-axial one.

The SCCM+RPA method, used in sections 4,5,6 to determine the structure and energies of the nuclear states near the yrast line, works well in the high-spin region where $J_x \approx J$ (see /3,4/). However, the question remains at what spins does this region begin. The first numerical calculations by this method, which have been done for low-lying positive parity states in ^{168}Er and ^{158}Dy in paper of Kvasil et al./9/, show that SCCM+RPA method provides quite good results for relatively small spins. The numerical calculations for negative parity states are to be prepared in the nearest future.

The authors express their gratitude to Professor I.N.Mikhailov for valuable discussions.

APPENDIX A: LINEAR MOMENTUM IN THE LAB. AND FIXED ROTATING SYSTEM

We suppose that the nucleus rotates in the lab. system around the stable axis. This axis of stable rotation can be chosen as axis x for the laboratory as well as for rotating system of coordinates. The origin point for both systems can be fixed in the center of mass of the nucleus in the ground state. It means that the total linear momentum in both systems is equal to zero for the nuclear states on the yrast line. When the intrinsic vibrational modes are excited in nucleus, the center of mass can be shifted from the origin point of both systems of coordinates and the total linear momentum in both systems can be non-zero. Since the rotation of the fixed system takes place around the x-axis, one can write

$$\begin{aligned} X' &= X, \\ Y' &= Y\cos\Omega t + Z\sin\Omega t, \\ Z' &= -Y\sin\Omega t + Z\cos\Omega t, \end{aligned} \tag{A1}$$

where (X',Y',Z') and (X,Y,Z) are the center of mass coordinates in the rotating and laboratory system, respectively. From the point of view of the classical physics the relation (A1) yields for the components of the total linear momentum:

$$\begin{aligned} P_{x'} &= P_x, \\ P_{y'} &= P_y\cos\Omega t + P_z\sin\Omega t + M\Omega Z', \\ P_{z'} &= -P_y\sin\Omega t + P_z\cos\Omega t - M\Omega Y'. \end{aligned} \tag{A2}$$

In the lab. system there is no strength acting on the nucleus centre of mass, and therefore, the classical Lagrangian for the center of mass has the form:

$$L = \frac{P_x^2+P_y^2+P_z^2}{2M} = \frac{P_{x'}^2}{2M} + \frac{(P_{y'}-M\Omega Z')^2}{2M} + \frac{(P_{z'}+M\Omega Y')^2}{2M}. \tag{A3}$$

The generalized momenta corresponding to coordinates in the rotating system are:

$$\begin{aligned} \tilde{P}_x &= \frac{\partial L}{\partial \dot{X}'} = P_x, \\ \tilde{P}_y &= \frac{\partial L}{\partial \dot{Y}'} = P_{y'} - M\Omega Z' = P_y\cos\Omega t + P_z\sin\Omega t, \\ \tilde{P}_z &= \frac{\partial L}{\partial \dot{Z}'} = P_{z'} + M\Omega Y' = -P_y\sin\Omega t + P_z\cos\Omega t. \end{aligned} \tag{A4}$$

From the Lagrangian equation one can get

$$\frac{d\tilde{P}_x}{dt} = \frac{\partial L}{\partial X'} = 0, \quad \frac{d\tilde{P}_y}{dt} = \frac{\partial L}{\partial Y'} = \Omega\tilde{P}_z, \quad \frac{d\tilde{P}_z}{dt} = \frac{\partial L}{\partial Z'} = -\Omega\tilde{P}_y. \tag{A5}$$

In quantum mechanics we must have in mind that $\tilde{P}_x$, $\tilde{P}_y$, $\tilde{P}_z$ are the canonical conjugated to coordinates X', Y', Z' (see (A4)). The transition from classical to quantum mechanics is:

$$\begin{aligned} \tilde{P}_x &\to -i\frac{\partial}{\partial X'} = e^{iJ_x\Omega t}\, P_x\, e^{-iJ_x\Omega t}, \\ P_y &\to -i\frac{\partial}{\partial Y'} = e^{iJ_x\Omega t}\, P_y\, e^{-iJ_x\Omega t}, \\ P_z &\to -i\frac{\partial}{\partial Z'} = e^{iJ_x\Omega t}\, P_z\, e^{-iJ_x\Omega t}. \end{aligned} \tag{A6}$$

The quantum mechanics generalization of relations (A5) gives then the conditions (2.5c) for Hamiltonian H'.

APPENDIX B: LINEAR BOSON PART OF SINGLE-PARTICLE OPERATORS

Every single-particle operator can be expressed in the form of an expansion in the two-quasiparticle bosons (see /3/). A particular form of each term in this expansion depends on the symmetry of a given single-particle operator under time reversion, complex conjugation, space reflection, and rotation of π around x-axis. The positive parity single-particle operators connected with $(\gamma_P=+1)$ part of the Hamiltonian H' are discussed in ref. /4/ therefore we give only the boson representation of the operators $Q_\mu[r]$, $\tilde{P}_\mu (\mu = x, y, z)$ and $\tilde{X}_\mu ([\tilde{X}_\mu, \tilde{P}_\mu] = i)$, which are connected with $(\gamma_P=-1)$ part of the Hamiltonian H'. Since the negative parity operators have no zero-order term in the boson expansion, we can write:

$$Q_\mu(r) = Q_\mu^{(1)}[r] + Q_\mu^{(2)}[r] + \ldots\ldots, \quad \tilde{P}_\mu = P_\mu^{(1)} + P_\mu^{(2)} + \ldots., \quad \tilde{X}_\mu = X_\mu^{(1)} + X_\mu^{(2)} + \ldots. \tag{B1}$$

The explicit expressions for the linear boson part of the operators P_μ, X_μ and $Q_\mu[r]$ are:

$$P_x^{(1)} = \sum \tilde{p}_{ij}^{x} (b_{i\bar{j}}^{+} + b_{i\bar{j}}) ,$$

$$P_y^{(1)} = \frac{1}{2} \sum_{ij} [\tilde{p}_{ij}^{y} (b_{ij}^{+} + b_{ij}) - \tilde{\tilde{p}}_{ij}^{y} (b_{\bar{i}\bar{j}}^{+} + b_{\bar{i}\bar{j}})],$$

$$P_z^{(1)} = \frac{i}{2} \sum_{ij} [\tilde{p}_{ij}^{z} (b_{ij} - b_{ij}^{+}) + \tilde{\tilde{p}}_{ij}^{z} (b_{\bar{i}\bar{j}} - b_{\bar{i}\bar{j}}^{+})], \qquad \text{(B2)}$$

$$X_x^{(1)} = -i \sum_{ij} \tilde{x}_{ij}^{x} (b_{i\bar{j}}^{+} - b_{i\bar{j}}) ,$$

$$X_y^{(1)} = \frac{-i}{2} \sum_{ij} [\tilde{x}_{ij}^{y} (b_{ij}^{+} + b_{ij}) - \tilde{\tilde{x}}_{ij}^{y} (b_{\bar{i}\bar{j}}^{+} - b_{\bar{i}\bar{j}})] ,$$

$$X_z^{(1)} = \frac{1}{2} \sum_{ij} [\tilde{x}_{ij}^{z} (b_{ij}^{+} + b_{ij}) + \tilde{\tilde{x}}_{ij}^{z} (b_{\bar{i}\bar{j}}^{+} + b_{\bar{i}\bar{j}})],$$

$$Q_x^{(1)}[r] = \sum_{ij} \tilde{q}_{ij}^{x}[r] (b_{i\bar{j}}^{+} + b_{i\bar{j}}) ,$$

$$Q_y^{(1)}[r] = \frac{1}{2} \sum_{ij} [\tilde{q}_{ij}^{y}[r](b_{ij}^{+} - b_{ij}) - \tilde{\tilde{q}}_{ij}^{y}[r](b_{\bar{i}\bar{j}}^{+} - b_{\bar{i}\bar{j}})],$$

$$Q_z^{(1)}[r] = \frac{-i}{2} \sum_{ij} [\tilde{q}_{ij}^{z}[r](b_{ij}^{+} + b_{ij}) + \tilde{\tilde{q}}_{ij}^{z}[r](b_{\bar{i}\bar{j}}^{+} + b_{\bar{i}\bar{j}})],$$

where the quasiparticle matrix elements are given by

$$\tilde{p}_{ij}^{x} = \sum_{k\ell} (P_x)_{k\ell}(u_{ki}\bar{v}_{\ell j} + v_{ki}\bar{u}_{\ell j}), \qquad \tilde{p}_{ij}^{y} = \sum_{k\ell} (P_y)_{k\ell}(u_{ki}v_{\ell j} - v_{ki}u_{\ell j}) ,$$

$$\tilde{\tilde{p}}_{ij}^{y} = \sum_{k\ell} (P_y)_{k\ell}(\bar{u}_{ki}\bar{v}_{\ell j} - \bar{v}_{ki}\bar{u}_{\ell j}), \qquad \tilde{p}_{ij}^{z} = \sum_{k\ell} (P_z)_{k\ell}(u_{ki}v_{\ell j} - v_{ki}u_{\ell j}),$$

$$\tilde{\tilde{p}}_{ij}^{z} = \sum_{k\ell} (P_z)_{k\ell}(\bar{u}_{ki}\bar{v}_{\ell j} - \bar{v}_{ki}\bar{u}_{\ell j}), \qquad \tilde{x}_{ij}^{x} = \sum_{k\ell} (X_x)_{k\ell}(v_{ki}\bar{u}_{\ell j} - u_{ki}\bar{v}_{\ell j}),$$

$$\tilde{x}_{ij}^{y} = \sum_{k\ell} (X_y)_{k\ell}(u_{ki}v_{\ell j} + u_{\ell j}v_{ki}) , \qquad \tilde{\tilde{x}}_{ij}^{y} = \sum_{k\ell} (X_y)_{k\ell}(u_{ki}v_{\ell j} + u_{\ell j}v_{ki}),$$

$$\tilde{x}_{ij}^{z} = \sum_{k\ell} (X_z)_{k\ell}(u_{ki}v_{\ell j} + u_{\ell j}v_{ki}), \qquad \tilde{\tilde{x}}_{ij}^{z} = \sum_{k\ell} (X_z)_{k\ell}(\bar{u}_{ki}\bar{v}_{\ell j} + \bar{u}_{\ell j}\bar{v}_{ki}),$$

$$\tilde{q}_{ij}^{x}[r] = \sum_{k\ell} (Q_x[r])_{k\ell}(u_{ki}\bar{v}_{\ell j} - v_{ki}\bar{u}_{\ell j}), \qquad \tilde{q}_{ij}^{y,z}[r] = \sum_{k\ell} (Q_{y,z}[r])_{k\ell}(u_{ki}v_{\ell j} + v_{ki}u_{\ell j}),$$

$$\tilde{\tilde{q}}_{ij}^{y,z}[r] = \sum_{k\ell} (Q_{y,z}[r])_{k\ell}(\bar{u}_{ki}\bar{v}_{\ell j} + \bar{v}_{ki}\bar{u}_{\ell j}) ,$$

where $(P_\mu)_{k\ell}$, $(X_\mu)_{k\ell}$ and $(Q_\mu[r])_{k\ell}$ are the matrix elements in the Goodman basis (see [3]) u_{ki} and v_{ki} are the matrices of the Bogolubov transformation from the single-particle operators to the single-quasiparticle operators (assignment as in ref.[3]).

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Цвек С., Квасил Я., Гориев Б. E4-83-684

ПСФ и восстановление трансляционной симметрии гамильтониана вращающегося ядра

Исследована трансляционная симметрия гамильтониана вращающегося ядра и ее влияние на решение уравнений движения в рамках приближения случайной фазы /ПСФ/. Из требования восстановления трансляционной симметрии вращающегося гамильтониана,нарушенной деформированным средним полем нуклонов, конструируются остаточные взаимодействия, которые в свою очередь ведут к появлению вибрационных состояний вблизи линии ираст ядра. С полученным таким образом полным гамильтонианом решаются уравнения движения в рамках ПСФ. В работе предложен способ выделения духовых мод среди ПСФ решений, связанных с трансляцией центра масс ядра в лабораторной системе. Найдены формулы для энергий и структуры вибрационных состояний отрицательной четности в окрестности линии ираст.

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Cwiok S., Kvasil J., Choriev B. E4-83-684

The RPA and the Restoration of Translation Symmetry of Rotating Nucleus Hamiltonian

The translation symmetry of Hamiltonian of rotating nucleus and its relation to the solutions of RPA equations are investigated. The requirement of restoration of translation symmetry of rotating nucleus Hamiltonian is used for construction of the residual interactions leading to the appearance of the odd parity rotational states near the nucleus yrast line. With these residual interactions the equations of motion in the framework of RPA are solved. The method of extraction of the spurious modes, connected with nucleus center of mass motion in lab.system, from the RPA solutions is proposed. The formulae for energies and structure of the odd parity vibrational states near the yrast line are found.

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

Communication of the Joint Institute for Nuclear Research. Dubna 1983