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THE EFFECT OF THE PAULI PRINCIPLE ON THE NONROTATIONAL STATES IN ODD-A DEFORMED NUCLEI

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1. INTRODUCTION

For the description of nonrotational states in deformed nuclei with an odd number of nucleons, one should take into account the quasiparticle-phonon interaction [1,2]. The energies and wave functions of the low-lying nonrotational states in deformed nuclei in the rare earth and actinide regions have been calculated in refs. [3,4] and other papers. The analysis of the experimental data and their comparison with the calculations of refs. [5-7] have shown [5,7] that on the whole the theory describes satisfactorily the experimental data. There are indications of a more complex structure of the low-lying states. In ref. [8] the rotational bands have been calculated taking into account the Coriolis interaction, and a good agreement has been obtained.

In recent years a good deal of attention is paid to the study of the two-phonon states in the doubly even nuclei and vibratioan states in odd-A deformed nuclei. The calculation of these states requires to take the Pauli principle into account consistently. The effect of the Pauli principle on the two-phonon components of the excited state wave functions has been studied in ref. [9]. It has been shown in ref. [10] that if the Pauli principle is taken into account, the admixtures to the one-phonon states become less and the energy centroids of the two-phonon collective states shift considerably towards higher energies. How the Pauli principle is taken into account in odd-A deformed nuclei is shown schematically in ref. [11]. The effect of the Pauli principle in odd-A spherical nuclei is studied in ref. [12].

Since the Pauli principle effects strongly the collective two-phonon states in doubly even deformed nuclei, it is expedient to investigate its effect on the vibrational states in odd-A deformed nuclei. Just this problem is the goal of the present paper.

2. EQUATIONS OF THE MODEL

The Hamiltonian of the quasiparticle-phonon nuclear model taking into account the secular equation, which defines the one-phonon energies, can be written as [13]

$$H_M = H_v + H_{vq} \, , \tag{1}$$

$$H_v = \sum_{q\sigma} \epsilon(q)\, a^+_{q\sigma} a_{q\sigma} - \frac{1}{4} \sum_{\substack{g=\lambda\mu i \\ g'=\lambda\mu i'}} \sum_{\sigma} \frac{Q^+_{g\sigma} Q_{g'\sigma}}{\kappa_0^{(\lambda\mu)} \sqrt{Y_g Y_{g'}}} \, , \tag{2}$$

$$H_{vq} = -\frac{\sqrt{2}}{4} \sum_{g\sigma} \sum_{qq'} \frac{v^{(-)}_{qq'}}{\sqrt{Y_g}} \{ (Q^+_{g\sigma} + Q_{g-\sigma}) (f^g(qq') B(qq';\mu-\sigma) + \bar{f}^g(qq') \bar{B}(qq';\mu-\sigma)) + \text{h.c.} \} . \tag{3}$$

We use the new definition of the phonon operator [14]

$$Q^+_{g\sigma} = \frac{1}{2} \sum \{ \psi^g_{qq'} A^+(qq';\mu\sigma) - \phi^g_{qq'} A(qq';\mu-\sigma) + \bar{\psi}^g_{qq'} \bar{A}^+(qq';\mu\sigma) - \bar{\phi}^g_{qq'} \bar{A}(qq';\mu-\sigma) \} , \tag{4}$$

which depends explicitly on the sign σ of the angular momentum projection into the symmetry axis, where $\sigma = \pm 1$,

$$A^+(qq';\mu\sigma) = \sum_{\sigma'} \delta_{\sigma'(K-K'),\sigma\mu}\, \sigma' a^+_{q\sigma'} a^+_{q'-\sigma'} \, ,$$

$$\bar{A}^+(qq';\mu\sigma) = \sum_{\sigma'} \delta_{\sigma'(K+K'),\sigma\mu}\, a^+_{q'\sigma'} a^+_{q\sigma'} \, , \tag{5}$$

$$B(qq';\mu\sigma) = \sum_{\sigma'} \delta_{\sigma'(K-K'),\sigma\mu}\, a^+_{q\sigma'} a_{q'\sigma'} \, ,$$

$$\bar{B}(qq',\mu\sigma) = \sum_{\sigma'} \delta_{\sigma'(K+K'),\sigma\mu}\, \sigma' a^+_{q\sigma'} a_{q'-\sigma'} \, . \tag{6}$$

Here $a^+_{q\sigma}$ is the quasiparticle creation operator, $q\sigma$ are the quantum numbers of the single-particle states, amid them there is a projection K of the angular momentum into the symmetry axis of the nucleus; $g = \lambda\mu i$, $\kappa_0^{(\lambda\mu)}$ is the constant of the isoscalar multipole forces, i is the root number of the secular equation for the one-phonon states, always $K>0$, $\mu \geq 0$; $\epsilon(q)$ is the quasiparticle energy, $u^{(\pm)}_{qq'} = u_q v_{q'} \pm u_{q'} v_q$, $v^{(\pm)}_{qq'} = u_q u_{q'} \pm v_q v_{q'}$; u_q and v_q are the coefficients of the canonical Bogolubov transformation, $f^g(qq')$ is the single-particle matrix element, $Y_g = Y_g(n) + Y_g(p)$ is the normalization function of the one-phonon state, specifying its collectivity. The form of Y_g and the other notation are given in ref. [13,14].

The wave function of the state with a fixed value K_0^π is

$$\Psi_n(K_0^\pi \sigma_0) = \{ \sum_{q_0} C^n_{q_0} a^+_{q_0\sigma_0} + \sum_{g_2\sigma_2} \sum_{q_3\sigma_3} D^n_{g_2 q_3} \delta_{\sigma_3 K_3 + \sigma_2\mu_2, \sigma_0 K_0} a^+_{q_3\sigma_3} Q^+_{g_2\sigma_2} \} \Psi_0 , \tag{7}$$

where Ψ_0 is the ground state wave function of a doubly even nucleus, $Q_{g\sigma}\Psi_0 = 0$, n is the state number with given K_0^π, $n = 1, 2, 3, \ldots$ In order to take into account the Pauli principle in the quasiparticle plus phonon components (7), we use the commutation relations and calculate

$$\sum_{\substack{\sigma_2\sigma_3\\ \sigma_2'\sigma_3'}} \delta_{\sigma_3 K_3+\sigma_2\mu_2,\sigma_0 K_0}\delta_{\sigma_3' K_3'+\sigma_2'\mu_2',\sigma_0 K_0}\langle Q_{g_2'\sigma_2'}\alpha_{q_3'\sigma_3'}\alpha^+_{q_3\sigma_3}Q^+_{g_2\sigma_2}\rangle = \tag{8}$$

$$=\delta_{g_2g_2'}\delta_{q_3q_3'}\{\delta_{K_3+\mu_2,K_0^+}\delta_{|K_3-\mu_2|,K_0}\}+S^{K_0}(g_2'q_3'|q_3g_2)+\bar{S}^{K_0}(g_2'q_3'|q_3g_2),$$

where

$$S^{K_0}(g'q'|qg)=-\sum_{q_2}\{\bar{\psi}^{g'}_{qq_2}\bar{\psi}^{g}_{q'q_2}\delta_{K+\mu,K_0}\delta_{K'+\mu',K_0}+\psi^{g'}_{qq_2}\psi^{g}_{q'q_2}\times \tag{9}$$

$$\times\sum_{\sigma\sigma'\sigma_3}\delta_{\sigma_3 K+\sigma\mu,\sigma_0 K_0}\delta_{\sigma_3 K'+\sigma'\mu',\sigma_0 K_0}\delta_{\sigma_3(K_2-K),\sigma'\mu'}\delta_{\sigma_3(K_2-K'),\sigma\mu}\},$$

$$\bar{S}^{K_0}(g'q'|qg)= \tag{9'}$$

$$=-\sum_{q_2}\{\psi^{g'}_{qq_2}\bar{\psi}^{g}_{q'q_2}\sum_{\sigma\sigma_3}\sigma_3\delta_{\sigma_3(K_2-K),\sigma\mu'}\delta_{\sigma_3(K-\mu),\sigma_0K_0}\delta_{\sigma\mu'+\sigma_3K',\sigma_0K_0}+$$

$$+\bar{\psi}^{g'}_{qq_2}\psi^{g}_{q'q_2}\sum_{\sigma\sigma_3}\sigma_3\delta_{\sigma_3(K_2-K'),\sigma\mu}\delta_{\sigma\mu+\sigma_3K,\sigma_0K_0}\delta_{\sigma_3(K'-\mu'),\sigma_0K_0}.$$

In the quasidiagonal with respect to S^{K_0}approximation, we have

$$-\sum_{\substack{\sigma\sigma_3\\ \sigma'\sigma_3'}}\delta_{\sigma_3K_3+\sigma\mu,\sigma_0K_0}\delta_{\sigma_3'K_3'+\sigma'\mu',\sigma_0K_0}\langle Q_{g'\sigma'}\alpha^+_{q_3\sigma_3}\alpha_{q_3'\sigma_3'}Q^+_{g\sigma}\rangle= \tag{10}$$

$$=\delta_{\lambda\mu,\lambda'\mu'}\ S^{K_0}(\lambda\mu i'q_3|q_3\lambda\mu i),$$

where

$$S^{K_0}(g'q_3|q_3g)=-\sum_{q_2}\{\bar{\psi}^{g'}_{q_3q_2}\bar{\psi}^{g}_{q_3q_2}\delta_{K_3+\mu,K_0}+ \tag{11}$$

$$+\psi^{g}_{q_3q_2}\psi^{g'}_{q_3q_2}(\delta_{K_3-K_2,\mu}\delta_{K_3+\mu,K_0}+\delta_{K_2-K_3,\mu}\delta_{|\mu-K_3|,K_0})\}.$$

In the diagonal with respect to S^{K_0} approximation the normalization condition of the wave function (7) has the form

$$\sum_{q_0}(C^n_{q_0})^2+\sum_{g_2q_3}(D^n_{g_2q_3})^2(1+S^{K_0}(g_2q_3|q_3g_2))=1. \tag{12}$$

We now calculate the average value of H_M by the wave function (7) and in the quasidiagonal with respect to $S^{K_0}(g_2'q_3|q_3g_2)$ approximation, we get

$$(\Psi^*_n(K_0^\pi\sigma_0)H_M\Psi_n(K_0^\pi\sigma_0))=\sum_{q_0}(C^n_{q_0})^2\epsilon(q_0)+$$

$$+\sum_{g_2q_3}(D^n_{g_2q_3})^2(1+S^{K_0}(g_2q_3|q_3g_2))[\epsilon(q_3)+\omega_g+\Delta^{K_0}(q_3g_2)]-$$

$$-\sum_{g_2q_3q_0}C^n_{q_0}D^n_{g_2q_3}\Gamma^{g_2}_{q_3q_0}(1+S^{K_0}(g_2q_3|q_3g_2)). \tag{13}$$

where

$$\Delta^{K_0}(q_3 g_2) = -\frac{1}{4\kappa_0^{(\lambda_2\mu_2)}} \sum_{i_4} \frac{S^{K_0}(\lambda_2\mu_2 i_4 q_3 | q_3 \lambda_2 \mu_2 i_2)}{\sqrt{Y_{\lambda_2\mu_2 i_4} Y_{\lambda_2\mu_2 i_2}}}, \quad (14)$$

$$\Gamma^{g_2}_{q_3 q_0} = \sqrt{2}\, \frac{v^{(-)}_{q_3 q_0}}{\sqrt{Y_{g_2}}} \{ f^{\lambda_2\mu_2}(q_0 q_3) \delta_{|K_3 - K_0|, \mu_2} + \sigma_0 f^{\lambda_2\mu_2}(q_0 q_3) \delta_{K_3 + K_0, \mu_2} \}, \quad (15)$$

Using the variational principle, we get, as in refs. [2,18], the following equations

$$\sum_{q_0'} C^n_{q_0'} \{ (\epsilon(q_0) - \eta_n) \delta_{q_0 q_0'} - \frac{1}{4} \sum_{g_2 q_3} \frac{\Gamma^{g_2}_{q_3 q_0} \Gamma^{g_2}_{q_3 q_0'} (1 + S^{K_0}(g_2 q_3 | q_3 g_2))}{\epsilon(q_3) + \omega_{g_2} + \Delta^{K_0}(q_3 g_2) - \eta_n} \} = 0, \quad (16)$$

$$D^n_{g_2 q_3} = \frac{1}{2} \frac{\sum_{q_0} C^n_{q_0} \Gamma^{g_2}_{q_3 q_0}}{\epsilon(q_3) + \omega_{g_2} + \Delta^{K_0}(q_3 g_2) - \eta_n}. \quad (17)$$

From the condition of existence of the nontrivial solution for the system (16), we get the secular equation for determining the state energies η_n of an odd nucleus

$$\theta(\eta_n) = \det \| (\epsilon(q_0) - \eta_n) \delta_{q_0 q_0'} - \frac{1}{4} \sum_{gq} \frac{\Gamma^{g}_{q q_0} \Gamma^{g}_{q q_0'} (1 + S^{K_0}(gq | qg))}{\epsilon(q) + \omega_g + \Delta^{K_0}(qg) - \eta_n} \| = 0. \quad (18)$$

From eqs. (12), (16) and (17) we derive explicit expressions for the functions $C^n_{q_0'}$ and D^n_{gq}. To this end, we choose some function $C^n_{q_0}$ and denote

$$\tilde{C}^n_{q_0'} = C^n_{q_0'} / C^n_{q_0}, \quad \tilde{D}^n_{gq} = D^n_{gq} / C^n_{q_0},$$

where $q_0 \neq q_0'$. We now rewrite eqs. (12), (16) and (17)

$$(C^n_{q_0})^2 \{ 1 + \sum_{q_0' \neq q_0} (\tilde{C}^n_{q_0'})^2 + \sum_{g_2 q_3} (\tilde{D}^n_{g_2 q_3})^2 (1 + S^{K_0}(g_2 q_3 | q_3 g_2)) \} = 1, \quad (19)$$

$$\tilde{C}^n_{q_0} (\epsilon(q_0') - \eta_n) - \frac{1}{2} \sum_{g_2 q_3} \tilde{D}^n_{g_2 q_3} \Gamma^{g_2}_{q_3 q_0'} (1 + S^{K_0}(g_2 q_3 | q_3 g_2)) = 0, \quad (20)$$

$$\tilde{D}^n_{g_2 q_3} = \frac{1}{2} \frac{\Gamma^{g_2}_{q_3 q_0} + \sum_{q_0' \neq q_0} \tilde{C}^n_{q_0'} \Gamma^{g_2}_{q_3 q_0'}}{\epsilon(q_3) + \omega_{g_2} + \Delta^{K_0}(q_3 g_2) - \eta_n}. \quad (21)$$

Hence, after simple transformations, we get

$$\tilde{C}^{n}_{q_0'} = \frac{\theta_{q_0}(q_0',\eta_n)}{\theta_{q_0}(\eta_n)}, \tag{22}$$

$$\tilde{D}^{n}_{g_2 q_3} = \frac{1}{2} \frac{\Gamma^{g_2}_{q_3 q_0} + \theta^{-1}_{q_0}(\eta_n) \sum_{q_0' \neq q_0} \theta_{q_0}(q_0',\eta_n)\Gamma^{g_2}_{q_3 q_0'}}{\epsilon(q_3) + \omega_{g_2} + \Delta^{K_0}(q_3 g_2) - \eta_n}, \tag{23}$$

$$(C^{n}_{q_0})^2 = \{1 + \theta^{-2}_{q_0}(\eta_n)[\sum_{q_0' \neq q_0} \theta^2_{q}(q_0',\eta_n) + \frac{1}{4}\sum_{g_2 q_3}(1 + S^{K_0}(g_2 q_3|q_3 g_2)) \times \tag{24}$$

$$- [\frac{\Gamma^{g_2}_{q_3 q_0'}\theta_{q_0}(\eta_n) + \sum_{q_0 = q_0} \theta_{q_0}(q_0'\eta_n)\Gamma^{g_2}_{q_3 q_0'}}{\epsilon(q_3) + \omega_{g_2} + \Delta^{K_0}(q_3 g_2) - \eta_n}]^2\}\}^{-1}, \tag{25}$$

$$\theta_{q_0}(\eta_n) = \det||(\epsilon(q_0'') - \eta_n)\delta_{q_0',q_0''} - V_{q_0'' q_0'}||,$$

$$V_{q_0 q_0''} = \frac{1}{4}\sum_{g_2 q_3} \frac{\Gamma^{g_2}_{q_3 q_0''}\Gamma^{g_2}_{q_3 q_0}(1 + S^{K_0}(g_2 q_3|q_3 g_2))}{\epsilon(q_3) + \omega_{g_2} + \Delta^{K_0}(q_3 g_2) - \eta_n}. \tag{26}$$

$\theta_{q_0}(q_0',\eta_n)$ is obtained by changing in (25) the columb q_0' by $V_{q_0'' q_0}$.

It is seen from (18) that if the Pauli principle is taken into account, there arises the factor $(1 + S^{K_0}(gq|qg))$, and the pole $\epsilon(q)+\omega_g$ shifts by the value of Δ^{K_0} (14). If the function $S^{K_0}(gq|qg)$ is taken to be equal to zero, eqs. (18)-(26) are reduced to formulae/13/ for which the Pauli principle was not taken into account.

3. SHIFT OF POLES OF THE QUASIPARTICLE PLUS PHONON IN THE SECULAR EQUATION

We now calculate the shift of poles of the quasiparticle plus phonon in the secular equation, which is caused by a consistent inclusion of the Pauli principle, and investigate the cases when the shifts of poles are large. The calculations are performed with the same model parameters as in ref./10/.

The largest two-quasiparticle components of the phonons in ^{166,168}Er, which may violate the Pauli principle in the quasiparticle plus phonon component of the wave function (7), are shown in Tabl.1. The shifts (14)and the values of the functions $S^{K_0}(gq|qg)$, entering the first term of the sum (14) are given in Tabl. 2. The calculated values of $S^{K_0}(gq|qg)$,

Table 1

Characteristics of the one-photon states

Nucleus	$\lambda\mu i$	ω_g, MэB	$Y_g(n)$	$Y_g(p)$	Phonon structure, %
^{168}Er	221	0.821	9	3	nn523↓521↓35% nn521↑521↓ 16% pp411↑411↓14% nn512↑521↓ 5%
	222	1.720	2050	12	nn512↑521↓94%
^{166}Er	221	0.786	15	3	nn523↓521↓39% nn521↑521↓ 30% pp411↑411↓ 9%
	311	1.700	$8\cdot10^4$	160	nn633↑523↓99%
	321	1.458	1860	230	nn633↑521↑62% nn642↑521↓ 12% pp523↑411↑ 7%

Table 2

The function $S^{K_0}(gq|qg)$ and the shift of the pole of the secular equation $\Delta^{K_0}(qg)$, arising in the equations for an odd-A nucleus taking into account the Pauli principle

Nucleus	K_0^{π}	$q \pm \lambda\mu i$	$S^{K_0}(gq\|qg)$	$\Delta^{K_0}(qg)$, MэB
^{169}Er	$3/2^-$	n521↓ −221	−0.48	3.0
	$5/2^-$	n521↓ +221	−0.23	2.2
	$1/2^-$	n523↓ −221	0	0
	$9/2^-$	n523↓ +221	−0.43	2.8
	$3/2^-$	n521↓ −222	−0.99	0.6
	$5/2^-$	n521↓ +222	−0.005	0.03
^{167}Er	$3/2^+$	n633↑ −221	−0.002	0.01
	$11/2^+$	n633↑ +221	−0.06	0.6
	$5/2^+$	n633↑ −311	−0.0003	$1\cdot10^{-4}$
	$9/2^+$	n633↑ +311	−0.99	0.08
	$3/2^+$	n633↑ −321	−0.003	0.02
	$11/2^+$	n633↑ +321	−0.65	1.7

and $\Delta^{K_0}(qg)$ indicate the degree of violation of the Pauli principle in the (gq) components at the given K_0.

The examples in the Tables clearly show the physical meaning of the corrections $S^{K_0}(gq|qg)$ and $\Delta^{K_0}(qg)$. We now consider, for instance, for ^{169}Er the one-quasiparticle state n 523↓ forming with the phonon $\lambda\mu i$ =221 the states with K_0=1/2 and 9/2. The quasiparticle n523↓ enters the neutron two-quasiparticle component nn523↓ 521↓ of the phonon $\lambda\mu i$ =221. The value of K_0=1/2 is obtained as a result of the projection summation $K_0 = K_{n523\downarrow} - \mu = K_{n523\downarrow} - (K_{n523\downarrow} - K_{n521\downarrow}) = \frac{1}{2}$ and the value of K_0=9/2 for $K_0 = K_{n523\downarrow} + \mu = K_{n523\downarrow} + (K_{n523\downarrow} - K_{n521\downarrow}) = \frac{9}{2}$. It is seen that in the first case the projections $K_{n523\downarrow}$ of the odd quasiparticle and of the quasiparticle in the phonon enter the summation with different signs, i.e., without violating the Pauli principle, and in the second case with the same signs, that is forbidden by the Pauli principle. Therefore, for $K_0 = \frac{1}{2}$ the corrections $S^{K_0=1/2}$ (g=211, q = n523↓| qg) and $\Delta^{K_0=1/2}$ (q=n523↓ , g=221) are equal to zero. For $K_0=\frac{9}{2}$, $S^{K_0=9/2}$ (g =221, q=n523 ↓| qg)=-0,43, and the pole $\epsilon(q=n523\downarrow) + \omega_{g=221}$ is shifted by the value of $\Delta^{K_0=9/2}$(q = =n523↓ g=221)=2,8 MeV. A large shift of the pole is due to a high collectivity of the state $\lambda\mu i$ =221 (see Tabl.1), i.e., to small values of $Y_g(n)$ and $Y_g(p)$ in (14).

Now we analyze the case of the one-quasiparticle state n521↓ forming with the phonon $\lambda\mu i$ =221 the state with K_0=3/2 and 5/2 and entering two main two-quasiparticle components of the phonon $\lambda\mu i$ =221. In the case of K_0=3/2 the projections are related by $K_0 = (K_{n523\downarrow} - K_{n521\downarrow}) - K_{n521\downarrow}$ =3/2 for the component nn 523↓ 521↓ and by $K_0 = (K_{n521\uparrow} + K_{n521\downarrow}) - K_{n521\downarrow}$ =3/2 for nn 521↑ 521↓ For K =5/2 we have $K_0 = (K_{n523\downarrow} - K_{n521\downarrow}) + K_{n521\downarrow}$ =5/2 and $K_0 = (K_{n521\uparrow} + K_{n521\downarrow}) + K_{n521\downarrow}$=5/2, respectively. It is seen that for K_0 =3/2 the Pauli principle is violated because of the component nn 523↓ 521↓ and for K_0=5/2 because of nn 521↑ 521↓ The corresponding values of S^{K_0} and Δ^{K_0} are shown in Tabl.2. For ^{169}Er the case for the quasiparticle n521↓ and phonon $\lambda\mu i$ =222 is also shown in Tabl.2. The comparison of the values of Δ^{K_0} (qg) for $\lambda\mu i$=221 and 222 shows that the shift of the pole Δ^{K_0} is to a great extent due to the phonon collectivity. The data for ^{167}Er are analysed in a similar way.

The analysis performed shows that if in the quasiparticle plus phonon components of the wave function (7) the Pauli principle is not violated or is slightly violated, then the pole shift equals zero or is small and the corresponding vibrational state has an energy close to the phonon $\omega_{\lambda\mu i}$ energy. This is the case with the γ-vabrational state $K_0^{\pi} = 3/2^+$ in

^{167}Er. If the Pauli principle is violated to a greater extent but not strongly, the pole shift is not large as for the state $K_0^{\pi}=11/2^+$ in ^{167}Er, and it may be responsible for the splitting of the γ-vibrational states ($q\pm2$). If the Pauli principle is strongly violated, the pole shift is large. Since at large energies the quasiparticle plus phonon state is fragmented strongly, in these cases the corresponding vibrational states should not be observed in odd-A deformed nuclei.

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Роль принципа Паули при описании неротационных состояний в нечетных деформированных ядрах

При использовании коммутационных соотношений между операторами квазичастиц и фононов получены уравнения, позволяющие корректно учесть принцип Паули при описании состояний нечетных деформированных ядер. Показано, что если в компоненте квазичастица плюс фонон принцип Паули не нарушен или нарушен слабо, то может существовать соответствующее вибрационное состояние в нечетном деформированном ядре.

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The Effect of the Pauli Principle on the Nonrotational States in Odd-A Deformed Nuclei

The commutation relations between the quasiparticle and phonon operators are used to obtain the equations allowing a correct inclusion of the Pauli principle for the description of the states of odd-A deformed nuclei. It is shown, that if in the quasiparticle plus phonon component the Pauli principle is not violated or is slightly violated, then a relevant vibrational state may exist in an odd-A deformed nucleus.

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

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