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THE STRUCTURE OF THE GROUND AND EXCITED STATES OF DEFORMED ODD-MASS NUCLEI IN THE ACTINIDE REGION

The interaction of quasi-particles with phonons in odd-mass deformed nuclei has been considered in a previous paper [1]. A secular equation for determining the nonrotational state energy has been derived. It has been shown that this interaction leads to the appearance of admixtures in states close to the single-particle ones and to the formation of collective nonrotational states and complex structure states. The structure of the ground and excited states of odd-mass nuclei in the region $153 \leq A \leq 187$ has been invetigated in ref.[2]. In the present note we give a part of the results obtained in invesitgating the structure of the excited states of odd-mass nuclei in the region $229 \leq A \leq 255$.

The secular equation determining the energies η_j of the ground and excited states of odd-mass deformed nuclei is of the form [1]

$$\varepsilon(\rho) - \eta_j - \frac{1}{4}\sum_{\lambda\mu i}\sum_{\nu} \frac{v_{\rho\nu}^2}{Y^i(\lambda\mu)} \frac{f^{\lambda\mu}(\rho\nu)^2 + \bar{f}^{\lambda\mu}(\rho\nu)^2}{\varepsilon(\nu) + \omega_i^{\lambda\mu} - \eta_j} = 0 \tag{1}$$

the collective state energies $\omega_i^{\lambda\mu}$ and the quantity $Y^i(\lambda\mu)$ being calculated in refs. [3,4,5], $\varepsilon(\nu) = \sqrt{C^2 + \{E(\nu) - \lambda\}^2}$ (C is the correlation function, λ is the chemical potential for an odd-mass nucleus), $v_{\rho\nu} = u_\rho u_\nu - v_\rho v_\nu$; $f^{\lambda\mu}(\rho\nu)$, $\bar{f}^{\lambda\mu}(\rho\nu)$ are the matrix elements of the multipole moment operator $(\lambda\mu)$. The summation over $\lambda\mu i$ is due to that one takes into account the interactions of quasi-particles with quadrupole $\lambda = 2$, $\mu = 0,2$ and octupole $\lambda = 3$, $\mu = 0,1,2$ phonons for the first two roots $i = 1,2$ of the secular equations for even-even nuclei, The wave functions and the Nilsson potential energies [6] are used in the calculations. The wave function describing the state with a given $K\pi$ is of the form:

$$\Psi(K\pi) = \Omega(K\pi)^+ \Psi_0 \tag{2}$$

$$\Omega(K\pi)^+ = \frac{1}{\sqrt{2}} C_\rho \left\{ \sum_\sigma \alpha_{\rho\sigma}^+ + \sum_{\lambda\mu i} \sum_{\nu\sigma} D_{\rho\nu\sigma}^{\lambda\mu i} \alpha_{\nu\sigma}^+ Q(\lambda\mu)^+ \right\}, \quad Q_i(\lambda\mu)\Psi_0 = 0 \tag{3}$$

where $Q_i(\lambda\mu)$ is the phonon operator of multipolarity $(\lambda\mu)$, $\alpha_{\nu\sigma}$ is the quasi-particle operator, $\sigma = \pm 1$, ρ stands for the average field levels with given $K\pi$'s and ν for the remaining levels.

$$C_\rho^{-2} = 1 + \frac{1}{4}\sum_{\lambda\mu i}\sum_{\nu} \frac{v_{\rho\nu}^2}{Y^i(\lambda\mu)} \frac{f^{\lambda\mu}(\rho\nu)^2 + \bar{f}^{\lambda\mu}(\rho\nu)^2}{(\varepsilon(\nu) + \omega_i^{\lambda\mu} - \eta_j)^2} . \tag{4}$$

$$D_{\rho\nu\sigma}^{\lambda\mu i} = \frac{1}{2} \frac{v_{\rho\nu}}{\sqrt{Y^i(\lambda\mu)}} \frac{f^{\lambda\mu}(\rho\nu) - \sigma \bar{f}^{\lambda\mu}(\rho\nu)}{\varepsilon(\nu) + \omega_i^{\lambda\mu} - \eta_j} . \tag{5}$$

The quantity C_ρ^2 determines the contribution of the one-quasi-particle state with a given ρ and $\frac{1}{2}C_\rho^2\sum_\sigma (D^{\lambda\mu i}_{\rho\nu\sigma})^2$ the contribution of the component with a quasi-particle in the ν state and a phonon with $\lambda\mu i$ to the considered state described by $\Psi(K\pi)$.

The investigations made in ref. [2] for nuclei in the region $151 \leq A \leq 187$ showed that the lowering of the energies η_j with respect to $\mathcal{E}(\rho)$ and to the first pole $\mathcal{E}(\nu)+\omega_1^{\lambda\mu}$ is mainly defined by the terms (1) with $\lambda=2, \mu=2, i=1$ and somewhat less strongly by the terms with $\lambda=3, \mu=0, i=1$. In some cases of importance are the terms in (1) with $\lambda=2, \mu=2, i=2$ and $\lambda=2, \mu=0, i=1$.

In the actinide region of great importance are phonons with $\lambda=2, \mu=0, i=1; \lambda=3, \mu=0$ $i=1$ and in some cases phonons with $\lambda=2, \mu=2; \lambda=3, \mu=1,2$ here i being unity. In the actinide region the role of beta-vibrational and octupole $(\mu=0)$ phonons essentially grew as compared with the rare-earth region. It should be noted that the terms in (1) with $\lambda>3$ and $i>2$ give a very small contribution since $Y^i(\lambda\mu)^{-1}$ tends to zero when the corresponding state of even-even nucleus approaches the two-quasi-particle one.

Each value of $K\pi$ has its own equation (1), the solutions for this equation are the energies $\eta_1, \eta_2, \ldots$ For the ground state of odd-mass nucleus $\eta_1(K_0\pi)$ assumes the smallest value and the energies of the excited states are the differences $\eta_j(K\pi)-\eta_1(K_0\pi)$.

In the cases when the interaction of quasi-particles with a beta-vibrational phonon played an important role then, instead of (1), eq. (13) was solved in ref.[2] in which the spurious state was excluded. However, this exclusion little changes the results of calculations. When in investigating the states with a given $K\pi$ there are several levels in the Nilsson scheme $\rho_1, \rho_2 \ldots \rho_n$ which have $\mathcal{E}(\rho_1), \mathcal{E}(\rho_2)\ldots\mathcal{E}(\rho_n)$ close to one another then, instead of (1) a complicated secular equation was solved. The general form of this equation which is a determinant of n-th order is given in [1] and the particular case $n=2$ is investigated in details in ref. [2]. The main role of this equation is the exclusion of false solutions for eq. (1) and the determination of the structure of relatively high states with a given $K\pi$.

The interaction of quasi-particles with phonons in the ground state of an even-even nucleus is taken into account, as in ref. [7], which leads to the appearance in (1) of an additional term without pole. The calculations performed showed that corrections due to the interaction of quasi-particles with phonons in the ground states of system with even number of nucleons enter the calculations errors and should be disregarded.

The secular equation (1) has no free parameter. The values of the poles in (1) are close to the two-quasi-particle state energies given in ref.[8]. The analysis of the solutions of eq. (1) shows that if η_1 is very close to $\mathcal{E}(\rho)$ then the state is close to the one-quasi-particle one. If η_1 noticeably differs from $\mathcal{E}(\rho)$ and the first pole $\mathcal{E}(\nu)+\omega_i^{\lambda\mu}$ then the structure of such a state is very complicated since, in addition to the one-quasi-particle state, many states with different quasi--particles and phonons contribute to the wave function. If η_1 is very close to the first pole of the secular equation then the state is collective.

We have calculated the energies of many levels for a large number of odd-mass nuclei in the region $229 \leq A \leq 255$. The wave functions have been found and the contribution of the one-quasi-particle states and of various components quasi-particle plus phonon has been calculated. As an example, in Table 1 we give the energies of the ^{237}Np levels up to 1 MeV and their structure (in per cent). The experimental values are taken from [6,9]. Most low-lying states of ^{237}Np are close to the one-quasi-particle ones, in some states the admixtures are very important and the 5/2 state of energy 0.9 MeV is close to the octupole one, and the 5/2 state of energy 1 MeV is almost purely the beta-vibrational one. From Table I it is seen that rather good description of the energy levels of ^{237}Np is obtained on which experimental data are available and the position of some additional states is predicted.

The interaction of quasi-particles with phonons relatively weakly affects the $K = 11/2$, $9/2$ states close to the one-quasi-particle states and somewhat more strongly the states with smaller K. This leads to different lowering with respect to the $\mathcal{E}(\rho)$ energies of these states. Therefore in a number of nuclei the calculated sequence of the excited states close to the one-quasi-particle ones differs from the sequence of the levels in the Nilsson scheme. The interaction of quasi-particles with phonons leads to a change of the state energy in different nuclei with identical odd values of N and Z. For instance, at $N=143$ the energy of the $K\pi = 1/2+$ state close to the $631\downarrow$ state in ^{235}U is equal to 0.08 KeV and in ^{237}Pu - 145 KeV. The calculations show that the energy of this state in ^{235}U is 10 KeV and in ^{237}Pu - 150 KeV. On the whole, the energies of the states close to the one-quasi-particle ones which are calculated taking into account the interaction of quasi-particles with phonons somewhat better agree with experimental data than the results obtained in the independent quasi-particle model [10].

The interaction of quasi-particles with phonons leads to the formation of collective nonrotational states in odd-mass nuclei and of complex structure states. Table 2 gives all the experimental data [9, 11-13] on such type states and the results of calculations. The $K\pi = 1/2-$ states of energy 685 KeV in ^{239}U the $K\pi = 1/2-$ states of energy

650 KeV in ^{235}U and the $K\pi = 1/2-$ states of energy 451 KeV in ^{239}Pu are to a large extent octupole ones. The $K\pi = 5/2-$ states in ^{237}Np of energy 721 KeV and in ^{239}Np of energy 766 KeV are beta-vibrational . It would be very interesting to determine experimentally the octupole $K\pi = 5/2-$ states in these nuclei which according to the calculations, are somewhat below than the beta-vibrational ones.

Thus, the account of the interaction of quasi-particles with phonons allowed to explain the position of all the experimentally found collective non-rotaional states and predict many states of such a type in odd-mass nuclei in the actinide region.

The properties of the collective states and the comlex structure ones as well as the admixture in states close to the one-quasi-particle states are revealed in the probabilities of the electrical E2 and E3 transitions, alpha- and beta-decay rates, the values of the spectroscopic factors in direct nuclear reactions, in the values of the decoupling parameters a for the $K\pi = 1/2$ states and so on. Let us consider the hindrance factors HF in alpha decays. If a $K_1\pi_1$ state close to the one-quasi-particle state ρ_1 alpha decays to the component quasi-particle ρ_1 plus phonon $Q_i(\lambda\mu)$ of the wave function with $K_2\pi_2$ and ρ_2 then the squared matrix element is of the form:

$$|M_\alpha(\rho_1 \to \rho_1 + Q_i(\lambda\mu))|^2 = C_{\rho_1}^2 \, C_{\rho_2}^2 \, \frac{1}{2}\sum_{\sigma} \left(D^{\lambda\mu i}_{\rho_2\rho_1\sigma}\right)^2 |M_{Q_i(\lambda\mu)}|^2 \tag{6}$$

where $M_{Q_i(\lambda\mu)}$ is the matrix element of the alpha transition from the ground to the collective $\lambda\mu i$ state of even-even nuclei the equation for which is given in ref.[3] The hindrance factor HF which defines the hindrance of the transition to the given state as compared to the transition of the same energy between the ground states of the corresponding even-even nuclei is of the form

$$HF(\rho_1 \to \rho_1 + Q_i(\lambda\mu)) = \frac{HF(Q_i(\lambda\mu))}{C_{\rho_1}^2 \, C_{\rho_2}^2 \, \frac{1}{2}\sum_{\sigma}\left(D^{\lambda\mu i}_{\rho_2\rho_1\sigma}\right)^2} \, , \tag{7}$$

where $HF(Q_i(\lambda\mu))$ is the hindrance factor for the alpha decay to one-phonon $\lambda\mu i$ state [3]. This transition is favourable and its hindrance as compared to the alpha decay to the collective one-phonon state of the even-even nucleus is due to the admixtures in the ρ_1 state of the parent nucleus and to a non-unity contribution of the state quasi-particle ρ_1 plus phonon in the wave function of the daughter nucleus. Such alpha transitions are somewhat hindered, as compared with the favourable ones, and noticeably enhanced as compared with the unfavourable ones, provided that $HF(Q_i(\lambda\mu))$ is not too large and the fraction of the considered state in $\rho_1 + Q_i(\lambda\mu)$ is not too small. The probabilities for the alpha transitions to other components (quasi-particle $\nu \neq \rho_1$ plus phonon) of the wave function with $K_2\pi_2$ and ρ_2 are about the same as those for the unfavorable alpha transitions between one-quasi particle states. A small value of (7) shows an essential admixture of the $\rho_1 + Q_i(\lambda\mu)$ state.

Table 2 gives the experimental data on the hindrance factors [9,11] and the calculated values. So, $HF(Q_1(20)) = 4$ for ^{234}U, according to [11], therefore $HF(631\downarrow \to 631\downarrow + Q_1(20)) \simeq 60$. Owing to the fact that there is no experimental data on $HF(Q_1(30))$ in ^{234}U, we take $HF(Q_1(30)) = 160$ i.e. like that in the alpha transitions $^{236}Pu \to {}^{232}U$ and $^{242}Cm \to {}^{238}Pu$ [11]. For the alpha decays $^{240}Pu \to {}^{236}U$ and $^{242}Pu \to {}^{238}U$ we take $HF(Q_1(20)) = 10$ i.e. like that in the alpha decay $^{242}Cm \to {}^{238}Pu$ and find $HF(523\downarrow \to 523\downarrow + Q_1(20)) = 10 \div 11$ in ^{237}Np and ^{239}Np. From Table 2 it is seen that the calculated values prove that the interpretation of the collective states in ^{235}U, ^{237}Np and ^{239}Np is correct. It should be noted that the hindrance factors on the octupole $K\pi = 5/2-$ states in ^{237}Np and ^{239}Np are essentially larger than on the beta-vibrational states; namely, this just results in the interpretation of the experimentally found states as beta-vibrational ones.

We investigate the effect of the interaction of quasi-particles with phonons on the decoupling parameters a. Taken the wave function $\Psi(K\pi)$ in the form (2) we get

$$a = C_\rho^2 \left\{ a^N_{\rho\rho} + \sum_{\nu\nu' i} a^N_{\nu\nu'} \left(D^{20i}_{\rho\nu+} D^{20i}_{\rho\nu'+} - D^{30i}_{\rho\nu+} D^{30i}_{\rho\nu'+} \right) \right\}, \qquad (8)$$

where $a^N_{\rho\rho}$, $a^N_{\nu\nu'}$ are the decoupling parameters a calculated with the Nilsson wave functions [6]. For states close to the pole with $\mu \neq 0$ the a is zero. The role of the second term in (8) in non-essential for the nuclei in the rare-earth region but essential for some nuclei in the actinide region. Table 3 gives the experimental values of the decoupling factor [9,11,12,14] and their calculated values, the values of a^N are also given, knowing C_ρ^2 it is easy to find $C_\rho^2 a^N$. To illustrate the role of the beta-vibtational terms and the octupole ones with $\mu = 0$ we give their contribution to the calculated a. Table 3 gives the values of a

for the collective states, the agreement between the calculated and experimental values for the $K\pi = 1/2^-$ state of energy 685 KeV in ^{239}U provides evidence for the correct description of the structure of this state while for the state of such a type in ^{239}Pu the situation is unclear. The calculated values of a for the states close to 631$\downarrow$ in the U and Pu isotopes are in their absolute value larger than the experimental ones. Perhaps, this is due to the defects of the Nilsson potential wave function. The calculated values of a for states close to 530$\uparrow$ are also larger than the experimental ones, the decrease of a as compared to a^N due to the multiplier C_ρ^2 being compensated by the addition from the octupole phonon. The account of the interaction of quasi-particles with phonons does not eliminate disagreement between the calculated and experimental values of a for states close to the one-quasi-particle ones, though it decreases this disagreement as compared to a^N.

We have calculated the properties of the ground and excited states for 30 nuclei in the region $229 \leq A \leq 255$, $20 \div 30$ states have been calculated for each nucleus. Thus we have accumulated a large amount of experimental material. Tables $1 \div 3$ give a small part of the results concerning the most interesting cases and the cases on which there are experimental data. The remaining material can be utilized as the amount of experimental data increases.

The aim of the present paper is to give a general picture of the excited states for many odd-mass nuclei. Therefore we have not performed a careful analysis of individual nuclei. Further one should analyse in detail the properties of the most interesting nuclei improving the Nilsson potential parameters, taking into account the Coriolis interaction and so on. With such an approach it is possible to obtain better agreement between theory and experiment and improve the predictions for the considered states.

It should be noted that in investigating the interaction of quasi-particles with phonons there is no free parameter, the quantities $\omega_i^{\lambda\mu}$ and $Y^i(\lambda\mu)$ are obtained in calculating the collective states of even nuclei. Therefore when the agreement between theory and experiment was insufficiently good in even nuclei this incorrectness is transferred to the description of odd-mass nuclei. On the whole, the general picture of the excited states of odd-mass nuclei is more complicated and the description more rough as compared to even nuclei. The investigations have shown that the structure of excited nonrotational states of deformed odd-mass nuclei is a rather various. If most low-lying states are close to the one-quasi-particle ones then, the energy increase, the number of states close to the collective ones and of complex structure states increases. The account of the interaction of quasi-particles with phonons has led to the improvement of the description of nuclear states close to the one-quasi-particle states as compared with the independent quasi-particle model and to a sufficiently correct description of the collective and the complex structure states. For further study of the structure of excited states of odd-mass deformed nuclei it is necessary to have a larger amount of experimental data on the states energies, beta and gamma-transition probabilities, spectroscopic factors in direct nuclear reactions and so on.

It should be noted that the position of the deformed odd-mass nucleus levels is to a large extent defined by the behaviour of one-particle average field levels. Therefore the accuracy of the calculation of different characteristics of odd-mass nuclei is resticted to a rough description of the energies and ths Nilsson potential wave function.

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Table 1

Energy and structure of the ground and excited states in ^{237}Np

$K\pi$	Energy (KeV) Exper.	Energy (KeV) Calcul.	Structure of the state $C_\rho^2 \; \frac{1}{2}\sum_\sigma (D^{\lambda\mu i}_{\rho\sigma s})^2$
5/2+	0	0	642↑ 93%; 642↑+Q(20) 3,2; 521↑+Q(31) 1.2
5/2−	59.6	140	523↓ 96%; 521↓+Q(22) 1%; 523↓+Q(20) 1%
1/2−	270	170	530↑ 82%; 530↑+Q(20) 5.9 %; 660↑+Q(30) 5.3%
1/2+	327	250	400↑ 79%; 400↑+Q(20) 10%; 402↓+Q(22) 4.4%
3/2+	(357)	260	651↑ 69%; 651↑+Q(20) 24%; 530↑+Q(31) 2.4%
1/2+		300	660↑ 53%; 530↑+Q(30) 20%; 660↑+Q(20) 19%
3/2−		425	532↓ 74%; 532↓+Q(20) 17%; 651↑+Q(30) 3.3%
3/2−	438	470	521↑ 84%; 642↑+Q(31) 5.2%; 521↑+Q(20) 3.6%
7/2−	−	580	633↑ 90%; 633↑+Q(20) 4.9%; 521↑+Q(32) 1.3%
11/2−	−	700	505↑ 69%; 505↑+Q(20) 30%
1/2−	−	800	541↓ 57%; 541↓+Q(20) 22%; 530↑+Q(20) 8.5%
5/2−	−	900	642↑+Q(30) 92%; 512↑ 5.8%; 633↑+Q(31) 0.5%
5/2−	721	1000	523↓+Q(20) 98%; 523↓ 1.1 %; 642↑+Q(30) 0.3%

Table 2

Energy, structure and forbiddenness factors for α – decay for the collective states and the complex structure states of deformed nuclei

Nuclei	$K\pi$	Energy (KeV)		Forbiddeness factor HF		Structure of the states
		Exp.	Calcul.	Exp.	Calcul.	
^{235}U	1/2–	650	660	75	300	631↓+Q(30) 52%; 761↓ 34%;761↓+Q(20)9%
^{235}U	1/2+	780	760	25	60	640↑59%; 640↑+Q(20) 24%;631↓+Q(20) 6.5%
^{237}U	1/2–	–	630	–	–	631↓+Q(30) 55%; 761↓ 33%;761↓+Q(20) 7.5%
^{239}U	1/2–	685	530	–	–	631↓+Q(30) 66%;761↓27%;761↓+Q(20)3.8%
^{239}Pu	1/2–	451	560	2460	–	631↓+Q(30) 81%; 761↓ 17%;
^{237}Np	5/2–	–	900	–	–	642↑+Q(30) 92%;512↑5.8%;633↓+Q(31)0.5%
^{237}Np	5/2–	721	1000	13	10	523↓+Q(20) 98%;523↑1.1%; 642↑+Q(30) 0.3%
^{237}Np	1/2+	327	250	2400	–	400↑ 79%; 400↑+Q(20) 10%; 402↓+Q(22) 4.4%
^{239}Np	5/2–	–	800	–	–	642↑+Q(30) 94%; 512↑ 5.2 %; 523↓+Q(20) 0.2%
^{239}Np	5/2–	666	930	24	11	523↓+Q(20) 97%; 642↑+Q(30) 0.5% 523↑0.4%
^{239}Np	1/2+	326	460	–	–	400↑ 77%; 530↑+Q(30) 13.4%; 400↑+Q(20) 8.4%

Table 3

Decoupling parameter

Nuclei	$K\pi$	ρ	Energy (Kev) Exp	Calcul	a Exp	Calcul	a^N	$C_\rho^2 \cdot a^N_{\nu\nu}(D^{201}_{\nu\nu+})^2$	$-C_\rho^2 a^N_{\nu\nu}(D^{301}_{\nu\nu+})^2$
^{235}U	1/2+	640	780	760	-	-0.6	-0.96	-0.23	0.21
^{235}U	1/2-	761	650	660	-	-0.84	-3.13	-0.28	0.50
^{239}U	1/2-	761	685	530	0.2	0.15	- 1.61	-0.06	0.64
^{239}Pu	1/2-	761	451	560	-0.4	0.25	- 1.61	-0.01	0.78
^{233}U	1/2+	631	399	250	-0.23	-0.65	- 0.89	-0.16	0.14
^{235}U	1/2	631	0.08	10	-0.30	-0.8	-0. 96	-0.16	0.08
^{237}U	1/2+	631	0	0	-0.44	-0.8	-0.96	-0.08	0.07
^{239}U	1/2+	631	133	70	-0.54	-0.84	-0.06	-0.01	0.08
^{237}Pu	1/2+	631	145	150	-0.4	- 0.86	-0.96	-0.08	0.06
^{239}Pu	1/2+	631	0	0	-0.58	-0.80	-0.96	-0.02	0.10
^{241}Pu	1/2+	631	163	100	-0.75	-0.9	-0.96	-0.01	0.07
^{251}Cf	1/2+	620	0	0	0.1	0.1	0.18	0	0.04
^{233}Pa	1/2-	530	0	0	-1.33	-2.5	-2.5	-0.01	-0.35
^{237}Np	1/2-	530	270	170	-1.65	-2.5	-2.5	-0.15	-0.34
^{239}Np	1/2-	530	267	160	-1.2	-2.4	-2.5	-0.01	-0.40
^{237}Np	1/2+	400	327	250	1.1	0.48	0.41	0.04	0.08
^{239}Np	1/2+	400	326	460	-	0.70	0.41	0.04	0.34