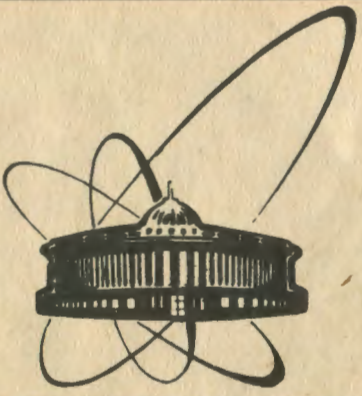

ОБЪЕДИНЕННЫЙ
ИНСТИТУТ
ЯДЕРНЫХ
ИССЛЕДОВАНИЙ
ДУБНА

E4-90-449

V.G.Soloviev, A.V.Sushkov, N.Yu.Shirikova

DESCRIPTION OF NONROTATIONAL STATES AND DISTRIBUTION OF Eλ STRENGTH IN DEFORMED NUCLEI IN THE ACTINIDE REGION

Submitted to "Ядерная физика"

1990

I. Introduction

The energies and wave functions of two-quasiparticle and one--phonon states of doubly even deformed nuclei were calculated in 1960–1975. Quadrupole and octupole states were calculated in the random phase approximation (RPA). The results of calculations are collected in refs.[1-3]. Good enough description of the available at that time experimental data was obtained and predictions were made most of which were later confirmed experimentally. In recent years, many new and more complete experimental data on the energies, $B(E\lambda)$-values and spectroscopic factors of one- and two-nucleon transfer reactions were obtained. Many experimental data are expected at the new generation of accelerators and detectors. Therefore, new, more exact and complete microscopic calculations of vibrational states of deformed nuclei are needed.

All subsequent calculations of nonrotational states of deformed nuclei were made in the quasiparticle-phonon nuclear model (QPNM) the basic assumptions of which are formulated in refs.[4-6]. In those calculations the wave functions of excited states were represented as sums of one-phonon and two-phonon terms. It was shown in [7] that due to the shift of two-phonon poles caused by the allowance of the Pauli principle, the energy centroids of collective two-phonon states exceed 2.5 MeV. At these energies the fragmentation of the two phonon strength should take place. Therefore, it was concluded in [7] that two-phonon collective states are absent in deformed nuclei. The calculations used not only quadrupole and octupole states. In ref.[8], the calculations of hexadecapole vibrational states showed that collective states with $K^{\pi} = 3^{+}$ should exist in the isotopes of Er, Yb, Hf and with $K^{\pi} = 4^{+}$ in the isotopes of Os. It was shown in ref.[9] that in some cases multipole interactions with $\lambda = 5 \div 9$ led to the mixing of two-quasiproton and two-quasineutron states in doubly even deformed nuclei and they should be taken into account.

Another stage of calculations is the inclusion of particle-particle (pp) interactions alongside with particle-hole (ph) ones[13-15]. The calculated quadrupole, octupole and hexadecapole states in ^{168}Er, ^{170,172,174}Yb and ^{178}Hf are in good agreement with experimental data. These calculations are limited by the states with $K^{\pi} \neq 0^{+}$. In ref. [15], equations for describing $K^{\pi} = 0^{+}$ states with inclusion of ph

and pp interactions were derived within the QPNM. From the condition of exclusion of spurious 0^+ states in the RPA there were derived equations for monopole and quadrupole pairing which have been studied in [16]. This allows one to include 0^+ states in the general scheme of calculations in the QPNM of all states of the electric type.

In the present paper, we give the transformed QPNM Hamiltonian for multipole ph and pp interactions which contains terms with $\mu=0$, i.e. including 0^+ states. Nonrotational states of some doubly even isotopes of Th, U and Pu are calculated and compared with experimental data. The E2 and E3 strength distribution in ^{238}U among the states with the excitation energy up to 6 MeV is given.

1. The Hamiltonian and basic equations of the QPNM

The initial QPNM Hamiltonian contains the average field of the neutron and proton systems in the form of a deformed axial-symmetric Woods – Saxon potential, monopole and quadrupole pairing and the effective ph and pp multipole interactions. To transform the Hamiltonian we use the canonical Bogolubov transformation

$$a_{q\sigma} = u_q \alpha_{q\sigma} + \sigma v_q \alpha^+_{q-\sigma}$$

and introduce the phonon operators $Q^+_{\lambda\mu i\sigma}$ and $Q_{\lambda\mu i\sigma}$ where

$$Q^+_{\lambda\mu i\sigma} = \frac{1}{2}\sum_{qq'}\left[\psi^{\lambda\mu i}_{qq'} A^+(qq';\mu\sigma) - \varphi^{\lambda\mu i}_{qq'} A(qq';\mu-\sigma)\right], \qquad (I)$$

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

$\alpha^+_{q\sigma}$ and $\alpha_{q\sigma}$ are the operators of creation and absorption of quasiparticles, $\psi^{\lambda\mu i}_{qq'}$ and $\varphi^{\lambda\mu i}_{qq'}$ are the direct and inverse amplitudes, $i=1,2,3,\ldots$ is the root number of the RPA secular equation. Single-particle states are specified by quantum numbers $q\sigma$ where $\sigma=\pm1$ and q includes K^π, where K is the angular momentum projection onto the symmetry axis of a nucleus and π is parity, and the asymptotic quantum numbers $N n_z \Lambda\uparrow$ at $K=\Lambda+1/2$ and $N n_z \Lambda\downarrow$ at $K=\Lambda-1/2$.

After transformations (see refs.[5,6]) the QPNM Hamiltonian is written in the form

$$H = \sum_{q\sigma} \varepsilon_q \alpha^+_{q\sigma} \alpha_{q\sigma} + H_{\upsilon} + H_{\upsilon q}\,, \qquad (2)$$

where the first two terms describe free quasiparticles and phonons, and the third describes the interaction of quasiparticles with phonons; ε_q is the quasiparticle energy with monopole and quadrupole

pairing [15,16]. Let us present the QPNM Hamiltonian in the explicit form, which has not been done before,

$$H_\upsilon = H_\upsilon^{00} + \sum_\lambda H_\upsilon^{\lambda 0} + \sum_{\substack{\lambda\mu \\ \mu\neq 0}} H_\upsilon^{\lambda\mu}, \tag{3}$$

$$H_\upsilon^{00} = -\frac{1}{2}\sum_{\tau i i'} G_\tau \left[d_{g\tau}^{i} d_{g\tau}^{i'} + d_{w\tau}^{i} d_{w\tau}^{i'} \right] Q_{20i}^{+} Q_{20i'}, \tag{4}$$

$$H_\upsilon^{\lambda 0} = -\sum_{ii'} W_{ii'}^{\lambda 0} Q_{\lambda 0 i}^{+} Q_{\lambda 0 i'}, \tag{5}$$

$$H_\upsilon^{\lambda\mu} = -\sum_{ii'\sigma} W_{ii'}^{\lambda\mu} Q_{\lambda\mu i\sigma}^{+} Q_{\lambda\mu i'\sigma}, \tag{6}$$

$$W_{ii'}^{\lambda 0} = \sum_\tau \Big\{ \sum_{p=\pm 1} (\varkappa_0^{\lambda 0} + p\varkappa_1^{\lambda 0}) D_\tau^{\lambda 0 i} D_{p\tau}^{\lambda 0 i'} + G^{\lambda 0} (D_{g\tau}^{\lambda 0 i} D_{g\tau}^{\lambda 0 i'} + D_{w\tau}^{\lambda 0 i} D_{w\tau}^{\lambda 0 i'}) \Big\}, \tag{5'}$$

$$W_{ii'}^{\lambda\mu} = \frac{1}{4}\sum_\tau \Big\{ \sum_{p=\pm 1} (\varkappa_0^{\lambda\mu} + p\varkappa_1^{\lambda\mu}) D_\tau^{\lambda\mu i} D_{p\tau}^{\lambda\mu i'} + G^{\lambda\mu} (D_{g\tau}^{\lambda\mu i} D_{g\tau}^{\lambda\mu i'} + D_{w\tau}^{\lambda\mu i} D_{w\tau}^{\lambda\mu i'}) \Big\} \tag{6'}$$

$$H_{\upsilon q} = H_{\upsilon q}^{00} + \sum_\lambda H_{\upsilon q}^{\lambda 0} + \sum_{\substack{\lambda\mu \\ \mu\neq 0}} H_{\upsilon q}^{\lambda\mu}, \tag{7}$$

$$H_{\upsilon q}^{00} = -\sum_{i\tau} G_\tau \sum_{qq'}^{\tau} (u_q^2 - \upsilon_q^2) u_{q'} \upsilon_{q'} \Big\{ (\psi_{qq'}^{20i} Q_{20i}^{+} + \varphi_{qq'}^{20i} Q_{20i}) \sum_\sigma \alpha_{q'\sigma}^{+} \alpha_{q\sigma} + \text{h.c.} \Big\}, \tag{8}$$

$$H_{\upsilon q}^{\lambda 0} = -\sum_{i\tau} \sum_{qq'}^{\tau} V_\tau^{\lambda 0 i}(qq') f^{\lambda 0}(qq') \Big\{ (Q_{\lambda 0 i}^{+} + Q_{\lambda 0 i}) B(qq'; \mu = 0) + \text{h.c.} \Big\}, \tag{9}$$

$$H_{\upsilon q}^{\lambda\mu} = -\frac{1}{2}\sum_{i\tau\sigma} \sum_{qq'}^{\tau} \Big\{ V_\tau^{\lambda\mu i}(qq') f^{\lambda\mu}(qq') Q_{\lambda\mu i\sigma}^{+} B(qq'; \mu - \sigma) + \text{h.c.} \Big\}, \tag{10}$$

$$V_{\tau}^{\lambda\mu i} = \frac{1}{2}\sum_{p=\pm1}(\varkappa_0^{\lambda\mu} + p\varkappa_1^{\lambda\mu})\,U_{qq'}^{(-)}\,D_{p\tau}^{\lambda\mu i} - \frac{1}{2}G^{\lambda\mu}U_{qq'}^{(+)}D_{q\tau}^{\lambda\mu i}\,. \qquad (10')$$

We use the following notation: $\varkappa_0^{\lambda\mu}$ and $\varkappa_1^{\lambda\mu}$ are the isoscalar and isovector constants of ph interaction of multipolarity λ with projection μ, $G^{\lambda\mu}$ is the constant of pp interaction, $f^{\lambda\mu}(qq') = \langle q \left| \frac{\partial V(r)}{\partial r} Y_{\lambda\mu}(\theta\varphi) \right| q' \rangle$ are single-particle matrix elements, where $V(r)$ is the central part of the Woods – Saxon potential. Summation over single-particle states of the neutron and proton systems is denoted by $\sum_{qq'}^{\tau}$ at $\tau=n$ and $\tau=p$, respectively,

$$d_{q\tau}^{i} = \sum_{q}^{\tau}\frac{E(q)-\lambda_\tau}{\varepsilon_q}\,g_{qq}^{20i}\,, \quad d_{w\tau}^{i} = \sum_{q}^{\tau} w_{qq}^{20i} \qquad (11)$$

$E(q)$ are single-particle energies, λ_τ are chemical potentials,

$$D_{\tau}^{\lambda\mu i} = \sum_{qq'}^{\tau} f^{\lambda\mu}(qq')\,U_{qq'}^{(+)}\,g_{qq'}^{\lambda\mu i}\,,$$

$$D_{q\tau}^{\lambda\mu i} = \sum_{qq'}^{\tau} f^{\lambda\mu}(qq')\,U_{qq'}^{(-)}\,g_{qq'}^{\lambda\mu i}\,, \qquad (11')$$

$$D_{w\tau}^{\lambda\mu i} = \sum_{qq'}^{\tau} f^{\lambda\mu}(qq')\,U_{qq'}^{(+)}\,w_{qq'}^{\lambda\mu i}\,,$$

$$B(qq';\mu\sigma) = \sum_{\sigma'}\delta_{\sigma'(K-K'),\sigma\mu}\,a_{q\sigma'}^{+}a_{q'\sigma'} \quad \text{or} \quad \sum_{\sigma'}\delta_{\sigma'(k+k'),\sigma\mu}\,\sigma' a_{q\sigma'}^{+}a_{q'-\sigma'}\,,$$

$$U_{qq'}^{(\pm)} = u_q v_{q'} \pm u_{q'}v_q\,, \quad V_{qq'}^{(\pm)} = u_q u_{q'} \pm v_q v_{q'}\,,$$

$$g_{qq'}^{\lambda\mu i} = \psi_{qq'}^{\lambda\mu i} + \varphi_{qq'}^{\lambda\mu i}\,, \quad w_{qq'}^{\lambda\mu i} = \psi_{qq'}^{\lambda\mu i} - \varphi_{qq'}^{\lambda\mu i}\,.$$

The one-phonon wave function can be written in the form

$$Q_{\lambda\mu i\sigma}^{+}\Psi_0\,, \qquad (12)$$

where Ψ_0 is the ground state wave function of a doubly even nucleus, defined as a phonon vacuum. Now, we find an average value

$$\sum_{q\sigma} \varepsilon_q \alpha^+_{q\sigma} \alpha_{q\sigma} + H_{v} \tag{13}$$

over the state (12) and using the variational principle we get the RPA equations. Their explicit form is given in [15] for $K^\pi = 0^+$ and in [13,14] for $K^\pi \neq 0^+$ states.

Nonrotational excited states of doubly even deformed nuclei will be calculated with the wave function

$$\Psi_\nu(K_0^{\pi_0}\sigma_0) = \Big\{ \sum_{i_0} R^\nu_{i_0} Q^+_{\lambda_0 K_0 i_0 \sigma_0} + \sum_{\substack{\lambda_1\mu_1 i_1\sigma_1 \\ \lambda_2\mu_2 i_2\sigma_2}} \frac{(1+\delta_{\lambda_1\mu_1 i_1, \lambda_2\mu_2 i_2})^{1/2}}{2\,[1+\delta_{K_0,0}(1-\delta_{\mu_1,0})]^{1/2}} \cdot \tag{14}$$

$$\cdot \delta_{\sigma_1\mu_1+\sigma_2\mu_2,\sigma_0 K_0} P^\nu_{\lambda_1\mu_1 i_1,\lambda_2\mu_2 i_2} Q^+_{\lambda_1\mu_1 i_1\sigma_1} Q^+_{\lambda_2\mu_2 i_2\sigma_2} \Big\} \Psi_0 ,$$

where $\nu = 1,2,3, \ldots$ is the state number with a given value of $K_0^{\pi_0}$. Using the variational principle we get equations for the energies E_ν and functions $R^\nu_{i_0}$ and $P^\nu_{\lambda_1\mu_1 i_1,\lambda_2\mu_2 i_2}$. Their explicit form is given in [6,14,15]. In the two-phonon terms of the wave function (14) the allowance is made for the Pauli principle.

2. Description of nonrotational states of nuclei in the actinide region

The calculations are made with single-particle energies and wave functions of the Woods-Saxon potential for the zones A=229 and 239. The parameters of the potential and equilibrium quadrupole and hexadecapole deformations are taken from [3]. The isovector ph constants are equal to $\varkappa_1^{\lambda\mu} = -1.5\,\varkappa_0^{\lambda\mu}$. The pp constants are equal to $G^{\lambda\mu} = 0.9\,\varkappa_0^{\lambda\mu}$. The constants of isoscalar ph interactions are taken in the range $\varkappa_0^{\lambda\mu} = 0.012 - 0.017$ fm^2 MeV^{-1} and $\varkappa_0^{44} = 0.020$ fm^{2}MeV^{-1}. The constants of the monopole pairing G_τ are calculated by pairing energies at fixed values of G^{20}. The energies of two-quasiparticle poles are calculated with inclusion of monopole and quadrupole pairing, blocking effect and the Gallagher-Moszkowski corrections.

The reduced probabilities of E2, E3 and E4 transitions from the ground state to the excited states $I^{\pi}K$ are calculated by the formulae (see /6,14/) with the effective charge $e^{eff}_{\tau} = 0.2$ and are given in single-particle units $B_{s.p.u.}(E\lambda) = \frac{2\lambda+1}{4\pi}\left(\frac{3}{\lambda+3}\right)^2 (0.12A^{1/3})^{2\lambda} e^2 10^{-24\lambda}$. Calculations of nonrotational states are made without the inclusion of the Coriolis interaction which should be taken into account for a detailed description of each nuclei. The calculated structure of a nonrotational state is given as a contribution (in per cent) of one-phonon $\lambda\mu i$ and two-phonon $\{\lambda_1\mu_1 i_1, \lambda_2\mu_2 i_2\}$ components to the normalisation of the wave function (14). The Pauli principle is taken into account in the contribution of two-phonon components. Then, we give (in per cent) several largest two-quasineutron nn and two-quasiproton pp components in the normalisation of the wave function (12) of the one-phonon state $\lambda\mu i$. The efficiency of calculations in the QPNM is demonstrated by ^{238}U and ^{240}Pu and nonrotational states of other nuclei in the actinide region.

The excited 0^+ states play an important role in the nuclear theory; they as though concentrate many difficulties of the theory. Description of several first 0^+ states of deformed nuclei in the RPA and QPNM with ph interactions and in other models /17/ cannot be thought to be satisfactory. In /18/, the energies and spectroscopic factors of (pt) and (tp) reactions with excitation of 0^+ states were calculated for some deformed nuclei of the rare-earth region. The calculations are made by an exact diagonalisation with allowance for effective pairing and neutron-proton quadrupole interactions. The observed enhancement of (tp) transitions to excited 0^+ states in some nuclei was reproduced by the calculation. It is a pity, the structures of the 0^+ states and B(E2) values were not presented.

As is shown in /13-14/, in describing $K^{\pi} \neq 0^+$ states in the QPNM, pp interactions do not influence greatly though somewhat improve the agreement with experimental data. In describing 0^+ states the role of pp interactions is considerable as with changing G^{20} the energies of several first poles of the RPA secular equation also change. At $G^{20} = 0.9\,\varkappa_0^{20}$ the B(E2) values for excitation of $I^{\pi}K_{\gamma} = 2^+0_1$ state and the energies of 0_2^+ and 0_3^+ states decrease and the structure of 0_1^+, 0_2^+ and 0_3^+ states changes in comparison with $G^{20} = 0$. The inclusion of pp interactions on the whole improves the description of 0^+ states.

The results of calculations of nonrotational states ^{238}U and ^{240}Pu and the experimental data /19-24/ are shown in Tables 1 and 2. Experimental data on the first three 0^+ states are insufficient and contradictory. Experimental data on 0^+ states in Tables 1 and 2 are

Table I

Nonrotational states in ^{238}U

K^{π}_{y}	Experiment E, MeV	Experiment B(Eλ) s.p.u.	Calculation E, MeV	Calculation B(Eλ) s.p.u.	Structure %					
0^{+}_{1}	(0.926)	0.4	0.95	0.8	201	96	{221,221}	1	201:	nn622↑ − 622↑ 36 nn624↓ − 624↓ 30 pp523↓ − 523↓ 10 nn631↓ − 631↓ 6
0^{+}_{2}	0.993	1.2	1.2	0.04	202	98			202:	nn624↓ − 624↓ 40 nn622↑ − 622↑ 36 nn631↓ − 631↓ 20
0^{+}_{3}	(1.482)	2	1.35	0.01	203	99				
0^{+}_{4}			1.6	0.08	204	92	{221,221}	1		
0^{+}	2.558	10^{-7}	3.2	$6\ 10^{-5}$						
2^{+}_{1}	1.06	3.0	1.1	4.0	221	98			221:	nn633↓ − 631↓ 20 nn622↑ − 631↓ 12
2^{+}_{2}	1.22	0.5	1.3	0.5	222	98			222:	nn622↑ − 631↓ 88 nn633↓ − 631↓ 4
2^{+}_{3}			1.9	0.1	223	64	224	25	223:	pp523↓ − 530↓ 77
0^{-}_{1}	0.68	25	0.63	18	301	98			301:	nn743↑ − 624↓ 19 nn752↑ − 662↑ 5
0^{-}_{2}			1.7	1.2	302	98			302:	nn743↑ − 624↓ 30 nn752↑ − 622↑ 6
1^{-}_{1}	0.93	8.1	0.97	11	311	98			311:	nn743↓ − 622↑ 62 pp642↑ − 521↑ 4
1^{-}_{2}	(1.113)		1.7	5	312	98			312:	nn743↓ − 622↑ 35 nn633↓ − 743↓ 14
2^{-}_{1}	1.129	6.4	1.1	6.7	321	96	{204,321}	1	321:	nn734↓ − 622↑ 25 pp642↑ − 530↑ 46
2^{-}_{2}			1.7	0.01	322	96	{204,321}	2	322:	nn734↑ − 612↑ 61 pp642↑ − 503↑ 38
3^{-}_{1}			1.6	0.2	331	98			331:	nn743↑ − 631↓ 96
3^{-}_{2}			1.9	2.7	332	92	{204,332}	3	332:	pp642↑ + 530↑ 58 nn752↑ − 631↓ 6
3^{+}_{1}	1.059		1.1	2.6	431	99			431:	nn631↓ + 622↑ 67 nn624↓ − 631↓ 19
3^{+}_{2}			1.5	0.4	432	98			432:	nn624↓ − 631↓ 68 nn622↑ + 631↓ 25
4^{+}_{1}			1.5	0.2	441	99			441:	nn624↓ + 631↓ 96
4^{+}_{2}			2.0	1.2	442	96			442:	nn613↓ + 631↓ 46 pp642↑ + 651↑ 12
5^{-}_{1}			1.6		551	100			551:	nn734↑ + 631↑ 93 pp642↑ + 523↓ 2
5^{-}_{2}			1.9		552	100			552:	nn752↑ + 622↑ 93 pp642↑ + 523↓ 5.8
5^{-}_{3}			2.0		553	100			553:	pp642↑ + 523↓ 87 nn752↑ + 662↑ 6.5

Table 2

Nonrotational states in ^{240}Pu

K^{π}_{ν}	Experiment E, MeV	Experiment B(Eλ) s.p.u.	Calculation E, MeV	Calculation B(Eλ) s.p.u.	Structure, %					
0^+_1	0.861		0.8	0.8	201	90			201:	nn622↑ – 622↑ 31
					{221,221}	0.5				nn624↓ – 624↓ 30
										pp523↓ – 523↓ 20
										nn631↓ – 631↓ 7
0^+_2	1.091		1.1	0.03	202	96			202:	nn622↓ – 622↑ 48
										nn624↓ – 624↓ 36
					203	2				nn631↓ – 631↓ 7
0^+_3	(1.525)		1.3	0.07	203	76	204	15	203:	nn631↓ – 631↓ 38
					202	3	201	2		pp523↓ – 523↓ 36
0^+_4			1.4	0.10	204	79			204:	pp523↓ – 523↓ 22
					203	18				nn743↓ – 743↑ 20
2^+_1	1.137	2.3	1.2	2.6	221	96			221:	nn622↑ – 631↓ 30
					{201, 221}	0.5				nn633↓ – 631↓ 21
										nn631↓ + 631↓ 7
										nn622↑ – 620↑ 6
2^+_2	(1.223)		1.3	1.0	222	96			222:	nn622↑ – 631↓ 69
					221	2				nn633↓ – 631↓ 10
0^-_1	0.597	17	0.6	13	301	94			301:	nn743↓ – 624↓ 17
										pp523↓ – 642↑ 14
										nn752↑ – 622↑ 4
0^-_2	(1.411)		1.2	0.3	302	88			302:	pp523↓ – 642↑ 34
										nn733↓ – 624↓ 15
0^-_3			1.7	2	303	98			303:	nn743↓ – 624↓ 18
										nn752↑ – 622↑ 17
1^-_1	0.938	9.6	0.94	7	311	96			311:	nn743↓ – 622↑ 65
										pp642↑ – 521↑ 17
1^-_2	(1.488)		1.4	2.2	312	94			312:	pp642↑ – 521↑ 58
										nn743↓ – 622↑ 31
2^-_1	1.241	10.8	1.24	8	321	98			321:	nn734↑ – 622↑ 45
										pp642↓ – 530↑ 12
										pp633↓ – 521↑ 6
2^-_2			1.9	1.5	322	88			322:	nn734↓ – 622↑ 50
					{221,301}	1				pp642↑ – 530↑ 25
3^-_1			1.5		331	99			331:	nn743↑ – 631↓ 96
3^+_1	1.036		1.0	1	431	99			431:	nn622↑ + 631↓ 90
4^+_1			1.35	0.3	441	99			441:	nn624↓ + 631↓ 90
										pp523↓ + 521↑ 6
5^-_1	1.308		1.2		551	100			551:	pp642↑ + 523↓ 100

shown for ^{238}U according to [19] and for ^{240}Pu according to [22]. In ^{238}U in ($\alpha,\alpha'\gamma$) and ($^{16}O,^{16}O'\gamma$) reactions [20] the $I^{\pi}K_{\nu}=2^{+}0_{1}$ state with energy 0.967 MeV is not observed. According to [25] 0_{1}^{+} and 0_{2}^{+} states in (tp) reaction are not excited. The 0_{3}^{+} state energy is not reliable. In some papers, the second 0^{+} state is treated as a beta-vibrational one. In [21] the first 0^{+} state is treated as a two--phonon gamma-vibrational state which we cannot agree with.

The first two 0^{+} states in ^{240}Pu are excited in (dp) and (pt) reactions. According to [24] the ratio $S_{\nu}(pt)/S_{g}(pt)$ of the spectroscopic factor for the 0_{ν}^{+} state to that for the ground state is equal to 0.15 for the 0_{1}^{+} state and 0.10 for 0_{2}^{+}. Excitation of 0_{1}^{+} and 0_{2}^{+} states in (dp) reaction means that in the one-phonon parts of their wave functions there should be the configuration nn 631↓– 631↓. The position with B(E2) values for $I^{\pi}K_{\nu}=2^{+}0_{1}$ and $2^{+}0_{2}$ states is indefinite as the rotational state $2^{+}0_{1}$ has the energy 0.900 MeV, and in (dd') reaction there excites the level with energy 0.939 MeV and B(E2) = 1.8 s.p.u. The energy of the rotational state $2^{+}0_{1}$ is close to the $2^{+}2_{1}$ energy and the corresponding B(E2) value is ascribed to the $2^{+}2_{1}$ state. The third 0_{3}^{+} state has the energy equal to 1.526 MeV according to [22] or the energy 1.410 MeV according to [23,26].

The calculations of 0^{+} states in ^{238}U and ^{240}Pu are made with the same constants $\varkappa_{0}^{20}=0.014$ fm^{2}MeV^{-1} and $G^{20}=0.9\varkappa_{0}^{20}$. According to the calculations, the wave functions of the first five 0^{+} states have dominating one-phonon components. The contribution of two-phonon components $\{221,221\}$ is equal or less than 1%. By changing the constants $\varkappa_{0}^{20}$ and G^{20} one can get small B(E2) values at state energies close to the experimental ones. However, one cannot obtain B(E2) value for the $2^{+}0_{2}$ state larger than for the $2^{+}0_{1}$ state, which takes place in ^{238}U. The energies of 0_{1}^{+} and 0_{2}^{+} states in ^{238}U are close, and for comparison of the theory with experiment it is insignificant which of them is below.

The calculated ratios $S_{\nu}(pt)/S_{g}(pt)$ for 0_{1}^{+} and 0_{2}^{+} states in ^{238}U and ^{240}Pu lie in the interval 0.01–0.02 which is one order of magnitude less than the experimental values [24] for ^{240}Pu. According to the calculations $S_{\nu}(pt)/S_{g}(pt)$ equals 0.3 for 0_{1}^{+} and 0.05 for 0_{2}^{+} states. For the third and higher 0^{+} states these ratios are small. In [18] and other papers it is shown that the increase of spectroscopic factors in (pt) and (tp) reactions with excitation of 0^{+} states is observed in nuclei which have a gap near the Fermi level in the position of energies of neutron single-particle states. For N = 146 in the position of single-particle levels there is no gap.

According to the calculations, in the wave functions of 0_1^+ and 0_2^+ states in ^{238}U and ^{240}Pu there is a considerable contribution of the configuration nn 631↓ – 631↓; therefore, these states should be excited in (dp) reactions, which is in agreement with experimental data /24/ for ^{240}Pu. According to /19/, at the energy 2.558 MeV in ^{238}U there is 0^+ state with $B(E2) = 10^{-7}$ s.p.u. for transition to 2^+0_g state. According to our calculations, there are states with energies 2.5 MeV and $B(E2) = 10^{-3}$ s.p.u. and 3.3 MeV and $B(E2) = 6\ 10^{-5}$ s.p.u. Our present calculations allow us to assert that with the pp interaction introduced the description of 0^+ states becomes better but not sufficient.

The data of Tables 1 and 2 show that there is good description of the energies and $B(E\lambda)$ values for $K^\pi = 2^+$ and first octupole states. Due to a relatively small pairing energy in ^{240}Pu the first poles with $K^\pi = 0^-, 1^-$ and 2^- lie very low, and therefore, too low calculated $B(E3)$ values are obtained in comparison with experimental ones. Experimental data on hexadecapole states in ^{238}U and ^{240}Pu are very scarce. According to the calculations, the wave functions of all nonrotational states up to 2 MeV have dominating one-phonon components.

Nonrotational states in ^{232}Th and ^{234}U are calculated. Description of experimental energies and $B(E\lambda)$ values is almost the same as in ^{238}U and ^{240}Pu. The calculated energies of $K_\gamma^\pi = 2_2^+$ states are 0.4 MeV higher than the experimental ones, and $B(E3)$ values for excitation of $I^\pi K_\gamma = 3^-0_1$ states are somewhat lower than the experimental ones. The energies and structure of hexadecapole states in ^{234}U are the following: $K_\gamma^\pi = 3_1^+$, $\mathcal{E}_1 = 1.5$ MeV, 431 99%, 431: nn 633↓ + 631↑ 40%; pp 523↓ + 530↑ 25%, nn 622↑ + 631↓ 16%; $K_\gamma^\pi = 3_2^+$, $\mathcal{E}_2 = 1.7$ MeV, 432 99%, 432: nn 633↓ + 631↓ 56%, pp 523↓ + 530↑ 22%, nn 622↑ + 631↓ 15%; $K_\gamma^\pi = 4_1^+$, $\mathcal{E}_1 = 1.8$ MeV, 441 96%, 441: nn 633↓ + 631↑ 25%, pp 642↑ + 651↑ 20%; $K_\gamma^\pi = 4_2^+$, $\mathcal{E}_2 = 1.9$ MeV, 443 88%, 444 5%; 443: nn 743↑ + 501↓ 80%. These results of calculations are in satisfactory agreement with the experimental data in /27/.

The energy centroids of two-phonon states are calculated. In these calculations with $G^{20} = 0.9\,\varkappa_0^{20}$ in comparison with the calculations /14/ with $G^{20} = 0$ the shift of two-phonon poles $\Delta\omega(2\nu_1, \lambda\mu_1)$ became considerably less, and therefore, the energy centroids of two-phonon states $\{2\nu_1, \lambda\mu_1\}$ decreased. Thus, in ^{238}U the shift $\Delta\omega(201, 221) = 2.4$ MeV at $G^{20} = 0$ and $\Delta\omega(201, 221) = 1.1$ MeV at $G^{20} = 0.9\,\varkappa_0^{20}$. According to the present calculations the largest shifts amount to 2.5 MeV for two-phonon poles $\{201, 301\}$. Other shifts $\Delta\omega(201, \lambda\mu_1)$ are less than 2.0 MeV; in the cases when phonon 201 does not participate the shifts are $\Delta\omega < 1.5$ MeV.

In the discussions, sometimes assertions are met that the $K_\gamma^\pi = 4_1^+$ state in ^{232}Th with the energy 1.414 MeV is the two-phonon state $\{221,221\}$. According to our calculations, in ^{232}Th the 4_1^+ state energy equals 1.42 MeV, the contribution of the one-phonon configuration 441 to the normalization of the wave function (14) equals 98% and of the two-phonon $\{221\ 221\}$ configuration equals 1%. In ^{234}U the 4_1^+ state has the following structure 441 96%, $\{201, 441\}$ 2%, $\{221,221\}$ 1%. The present calculations confirm the assertion, made in /7/, about the absence of two-phonon collective states in deformed nuclei.

3. Distribution of E2 and E3 strength in ^{238}U

It is assumed that the first quadrupole beta and gamma vibrational states are collective and other collective states form giant isoscalar and isovector quadrupole resonances. In the case of octupoles, there are first octupole collective states and other collective states are those forming low-lying and high-lying isoscalar and isovector octupole giant resonances. In some nuclei, low-lying collective hexadecapole states are observed. The collectivity of the first quadrupole and octupole states and its absence in higher lying states up to the giant resonances underlie phenomenological models including the interacting boson model (IBM).

Now, let us study the distribution of $E\lambda$ strength among the low-lying states with an energy up to 5 MeV. In the calculations we neglect the relation between the rotational and vibrational motions. Investigation of $E\lambda$ strength distribution became possible after the experiments in /28/ on inelastic scattering of α particles. It was shown that in ^{168}Er in accordance with experimental data /28/ among the $K^\pi = 3^-$ states most collective is not the first 3_1^- but fourth 3_4^- state. In ^{238}U according to our calculations the most collective are the second 3_2^- and fourth 3_4^- states with B(E3) equal to 2.7 and 2.5 s.p.u. whereas for the first 3_1^- state B(E3)= 0.2 s.p.u. Similarly, for the second 4_2^+ state B(E4) is six times larger than for the first 4_1^+ state. For other K values in ^{238}U the first quadrupole, octupole and hexadecapole states are the most collective among the low-lying states.

Let us demonstrate the E2 and E3 strength distribution among the low-lying vibrational states using ^{238}U as an example. To study the E2 strength distribution in the RPA we calculated B(E2; i) values for $K_i^\pi = 2_i^+$ and 0_i^+ states. The low-lying 1^+ states are close to two-quasiparticle ones and we neglect their contribution . The results of calculations are shown in Fig.1 as the sum $\sum_i$ B(E2;i) containing the parts corresponding to $K^\pi = 2^+$ and 0^+ . About 2/3 of the

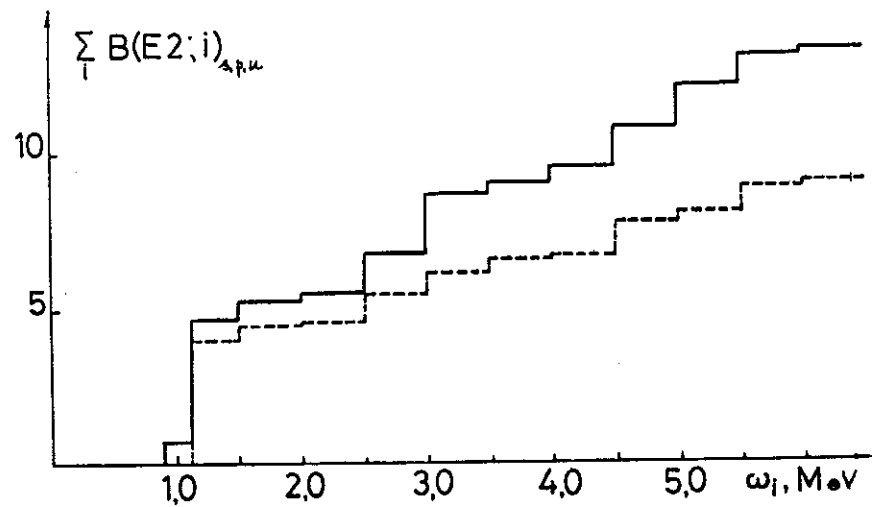

Fig.1. Distribution of E2 strength in ^{238}U represented as a total sum of B(E2,i) values over the energy intervals of 0.5 MeV. Notation: the dashed histogram is the contribution to $\sum_i B(E2;K=20)$ of the states with $K^\pi=2^+$, and the solid line is the contribution of the states with $K^\pi = 0^+$ and 2^+.

total sum $\sum_i B(E2;i)$ is concentrated in the $K^\pi=2^+$ state. It is seen from the figure that up to 5.5 MeV the total sum increases and then goes out to the plateau. For the states with energies more than 5.5 MeV up to the isoscalar quadrupole resonance B(E2;i) values are very small. The contribution of the gamma-vibrational $K_i^\pi = 2_1^+$ state to the sum $\sum_i B(E2;K=2,i)$ equals 45% and the contribution of the beta-vibrational state $K_i^\pi = 0_1^+$ to the sum $\sum_i B(E2;K=0,i)$ equals 15%.

For all the states with the energy up to 6.5 MeV the contribution of the sum $\sum_i \omega_{22i} B(E2;K=2,i) + \sum \omega_{20i} B(E2;K=0,i)$ to the energy weighed isoscalar sum rule is 15.5%, the contribution of beta-and gamma-vibrational states amounts to 2.2%. It follows from the investigations performed that a considerable part of the E2 strength is concentrated in states lying above beta- and gamma-vibrational states. The distribution of E2 strength, shown in fig.1, is almost unchanged in calculations with the wave function (14); thus, the shift of strength from one energy interval to another is not large. Note that phenomenological models inclusing IBM do not reproduce the part of the E2 strength which is concentrated in states lying above beta- and gamma-vibrational states.

The E3 strength distribution in ^{238}U calculated in the RPA in the form of B(E3,i) values for each state i with $K^\pi=0^-,1^-,2^-,3^-$ is shown in fig.2. It is seen from this figure that apart from the first there is a large number of collective octupole states. With increasing excitation energy up to 5.5 MeV B(E3) values decrease; above 6 MeV up to the isoscalar high-lying-energy octupole resonance they are very small. In the region of 3.5–4.0 MeV there is a local maximum which can hardly be treated as a isoscalar low-energy octupole resonance. According to the calculations, there is no gap in the distribution of B(E3) values between the first and higher lying octupole states. All octupole states with the energy up to 5.0 – 5.5 MeV can be thought to form the isoscalar low-energy octupole resonance. For ^{238}U the contribution of these states to the isoscalar energy weighed sum rule amounts to 28%. It is close to the value obtained in /29/ for ^{232}Th.

As is seen from fig.2, at energies 3.5, 4.8 and even 5.4 MeV there are collective octupole states with large B(E3) values. The calculations of $K^\pi = 0^-,1^-,2^-$ and 3^- states with the wave function (14) have shown that a large number of states lying in the interval 2–5 MeV have complex structure. In many cases, states with small B(E3) values are not two-quasiparticle and their wave functions contain a great number of two-quasiparticle and two-phonon components.

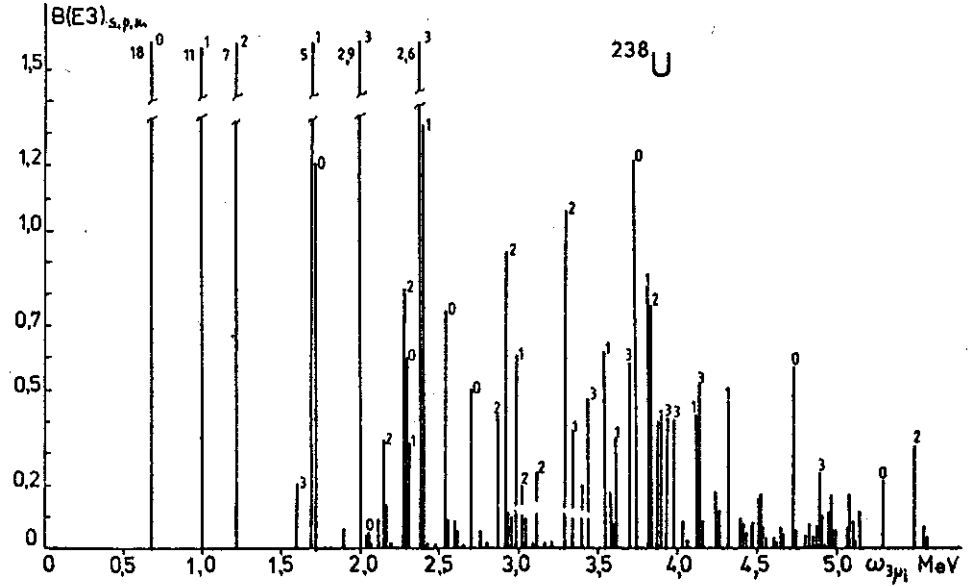

Fig. 2. Distribution of E3 strength among the low-lying states in ^{238}U. Notation: B(E3) are given in single-particle units, large B(E3) values are given on the left from the straight lines and K=0,1,2,3 are given on the right for B(E3)$>$ 0.2 s.p.u.

Some results are given in Table 3 where the same notation as in Table 1 is used. In the wave functions of many states the one-phonon component dominates; in some states there are large enough two-phonon components. In spite of the absence of three-phonon terms in the wave function (14) which influence the fragmentation of two-phonon collective states, the strength of each two-phonon collective state is distributed over many nuclear levels. Apparently, experimental study of the two-phonon strength distribution seems to be very timely.

The study of the E2 and E3 strength distribution in ^{238}U has shown that the major part of the strength concentrated in low-lying states lies above beta- and gamma-vibrational and first octupole states. Experimental study of the $E\lambda$ -strength distribution in doubly even deformed nuclei is obviously of particular scientific interest.

Conclusion

The performed investigations have shown that the wave functions of nonrotational states with energies up to 2 MeV have dominating one-phonon components. The total contribution of two-phonon components does not exceed 10%. Qualitatively correct description of the available experimental data is obtained which is exemplified by ^{238}U and ^{240}Pu. Some predictions are made. It is shown that apart from the first quadrupole and octupole states there should be a great number of vibrational states in the energy interval 2-4 MeV with B(E2) and B(E3) values equal to 0.5 - 1.0 s.p.u. Collective hexadecapole states are to be observed too. Therefore, experimental study of the $E\lambda$ -strength distribution among the states up to the excitation energy of 5 MeV is necessary. Note that there are no considerable difficulties in describing nonrotational states of nuclei in the actinide region with the number of neutrons $N > 140$ in comparison with nuclei of the rare-earth region.

Characteristics of excited states of deformed nuclei were calculated within the IBM (see for instance [11,12]), though mainly the rare-earth nuclei were used. In the actinide region, for instance in [30,31] there were calculated rotational bands based on the beta and gamma-vibrational states. As the low-lying 0^+ states lie outside the IBM space, as in shown in [32], the calculation in the IBM of nuclei in the actinide region presents difficulties. The mixing of configurations leads to a considerable increase in the number of parameters. Complication in the actinide region is due to low-lying states of negative parity, for the description of which apart from the f boson one should introduce the p boson [33]. Of certain difficulty of the description in the IBM is the absence of 0^+ states constructed of two f bosons with $K^\pi = 0^-$.

Table 3

Structure the selected states in ^{238}U

K^{π}_{ν}	$\mathcal{E}_{\nu}$ MeV	B(E3) s.p.u.	Structure, %
0^{-}_{16}	3.73	1.50	30 16 24; 201, 302 21; {331,434} 12; {205,301} 8 30 16: pp521↑ – 651↑ 8; pp512↑ – 642↑ 7; nn633↓ – 503↓ 4
0^{-}_{28}	4.80	0.29	30 28 82; 334,434 8 {201,301} 4 30 28: nn620↑–501↓ 36; nn761↑ – 622↓ 3
1^{-}_{27}	3.49	10^{-4}	31 27 97; {201,312} 1 31 27: nn501↓ – 631↑ 98
2^{-}_{33}	4.14	9.03	32 33 92 {321,434} 3; {322,442} 2 32 33: nn615↓ – 502↑ 42; pp514↓ – 402↓ 41
3^{-}_{4}	2.40	2.5	334 86; {221,311} 5 334: pp642↑+ 530↑ 29; nn725↑ – 622↑ 23; pp633↓ – 530↑ 10
3^{-}_{9}	2.91	0.1	339: nn725↑ – 622↑ 48; nn734↑ – 631↑ 44; pp633↓ – 530↑ 6 339 94; {201, 332} 2

In /34/ new phonon operators consisting of the electric and magnetic parts were introduced and the QPNM equations for the finite rank separable multipole and spin-multipole interactions were derived. The basis is created for describing the distribution of $M\lambda$ – strength and $E\lambda$ and $M\lambda$ transitions between excited states of deformed nuclei. Thus, it is shown that the QPNM can serve as a basis for further calculations of many characteristics of excited states of atomic nuclei.

We hope that these calculations will be useful for experiments at a new generation of accelerators and detectors.

References

1. Soloviev V.G. Theory of complex nuclei, Moscow, Nauka, 1971.
2. Grigoriev E.P., Soloviev V.G. Structure of even deformed nuclei, Moscow, Nauka, 1974.
3. Ivanova S.P., Komov A.L., Malov L.A., Soloviev V.G. Part.Nucl., 1976, v.7, p.450.
4. Soloviev V.G. Part.Nucl., 1978, v.9, p.580.
5. Soloviev V.G. Prog .Part.Nucl.Phys., 1987. v.19, p.107.
6. Soloviev V.G. Theory of atomic nuclei. Quasiparticle and Phonons, Moscow, Energoatomizdat, 1989.
7. Soloviev V.G., Shirikova N.Yu. Z. Phys.A – Atoms and Nuclei, 1981, v.301, p.263.
   Soloviev V.G., Shirikova N.Yu. Yad.Fyz., 1982, v.36, p.1376.
8. Nesterenko V.O., Soloviev V.G., Sushkov A.V., Shirikova N.Yu. Yad.Fyz., 1986, v.44, p.1443.
9. Soloviev V.G., Sushkov A.V. J.Phys. G. Nuc .Part.Phys., 1990, v.16, p.L57.
10. Soloviev V.G. Pisma v JETP, 1984, v.40, p.399;
    Yad.Fyz., 1988, v.47, p.332.
    Soloviev V.G. Z.Phys.A – Atomic Nuclei, 1986. v.324, p.393.
11. Yoshinaga N., Akiyata Y., Arima A. Phys.Rev., 1988, v.C38, p.419.
12. Barfield A.F., Barrett B.R., Wood J.L., Scholten O. Ann.Phys., 1988, v.182, p.344.
13. Soloviev V.G., Shirikova N.Yu. Izv. AN SSSR ser.fyz., 1988, v.52, p.2005.
14. Soloviev V.G., Shirikova N.Yu. Z. Phys. – Atomic Nuclei, 1989. v.334, p.149; Izv. AN SSSR, ser.fyz., 1990, v.54, p.818.
15. Soloviev V.G. Z.Phys. – Atomic Nuclei, 1989, v.334, p.143.
16. Karadjov D., Soloviev V.G., Sushkov A.V. Izv.AN SSSR, ser.fyz., 1989, v.53, p.2150.
17. Kusmenko N.K., Mikhailov V.M. Izv. AN SSSR, ser.fyz., 1979, v.43, p.2082.
    Abrosimov V.I., Strutinski V.M. Yad.Fyz., 1979, v.29, p.355.
    Mitropolsky I.A. J.Phys. G: Nucl. Phys., 1981, v.7, p.921.
18. Shinab A.A., Rasmussen J.O., Stoyer M. Workshop on Microscopic Models in Nuclear Structure Physics ed. by M.W.Guidry J.H.Hamilton et al. Singapore, World Scientific, 1989, p.282.
19. Shurshikov E.N. Nucl. Data Sheets, 1988, v.54, p.603.
20. Alessi J.G., Saladin J.X., Baktash C., Humanic T. Phys.Rev., 1981, v. C23, p.79.

21. Chan D.W., Sheldon E. Phys.Rev., 1982, v.C28, p.861.
22. Shurshikov E.N., Sergeev V.O., Jaborov Yu.F. Nucl. Data Sheets, 1984, v.43, p.245.
23. Sheldon E., Chan D.W. J. Phys.G: Nucl. Phys., 1987, v.13, p.227.
24. Friedman A., Kator K. Phys.Rev.Lett., 1973, v.30, p.102.
25. Casten R.F., Flynn E.R., Garritt J.O. e.a. Phys.Lett., 1972, v.40B, p.333.
26. Schmorak M.R., Bemis C.E., Zender M. et al. Phys.Rev.Lett., 1970, v.24, p.1507.
27. Ardisson C., Dalmasso J., Ardiccon G. Phys.Rev., 1986, v.C33, p.2132.
28. Govil I.M., Fulbright H.W. et al. Phys.Rev., 1986, v.C33, p.793; 1987, v.C36, p.1442.
29. Malov L.A., Soloviev V.G. Part.Nucl., 1980, v.11, p.301.
30. Golovkov N.A., Dzhelepov B.S., Ivanov A.B., Mikhailova M.A. Izv. AN SSSR, ser.fyz., 1986, v.50, p.2.
31. Zhang M., Valliores M., Gilmore R. et al. Phys.Rev., 1985, v.C32, p.1076.
32. Wood J.L. Nucl. Phys., 1984, C.A401, p.43C.
33. Engel J., Iachello F. Nucl. Phys., 1987, v.A472, p.61.
34. Soloviev V.G. Preprint JINR E4-90-119, Dubna, 1990.

Received by Publishing Department
on June 14, 1990.