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COMPLICATION OF THE NUCLEAR STATE STRUCTURE WITH INCREASING EXCITATION ENERGY

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COMPLICATION OF THE NUCLEAR STATE STRUCTURE WITH INCREASING EXCITATION ENERGY

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I. The generally accepted description of the nuclear spectra is based on the division of the nuclear motions into three types: rotational, vibrational and quasiparticle. Using the semi-microscopic approach a basis has been created for a unified description of both the quasiparticle and vibrational motions. Some premises have also been created for including into this scheme the rotational motion.

The description of the excited states of complex nuclei in terms of the quasiparticles and phonons is widely used. This description is rather good for doubly magic nucleus and the neighbouring nuclei and good enough for the ground and low-lying states of strongly deformed nuclei. As going away from the doubly magic nuclei and with increasing excitation energy the structure of the levels becomes more complicated due to admixtures to the quasiparticle and phonon components.

There are two main causes leading to the appearance of admixtures to the quasiparticle and phonon states. This is the coupling of the internal motion with the rotation of the nucleus as a whole and the interaction of quasiparticles with phonons. A large amount of new experimental data has allowed to conclude that the low-lying states of spherical and deformed nuclei turned out to be essentially more complicated compared to the conceptions which have recently been formulated. Therefore, to continue the study of the structure of low-lying nuclear states and go further to higher excitations it is necessary to combine different methods: alpha-, beta- and gamma-ray spectroscopy (including $(n\gamma)$ reactions) with direct nuclear reactions.

2. We present here the basic formulas of the theory in which the interaction of quasiparticles with vibrational phonons of the corresponding even-even nuclei is taken into account[1,2].

The wave function of an odd-A nucleus describing the states with given K^{π} is written in the form

$$\Psi_i(K^{\pi}) = C^{i}_{\rho}\frac{1}{\sqrt{2}}\left\{\sum_{\sigma}\alpha^{+}_{\rho\sigma} + \sum_{\lambda\mu j}\sum_{s\sigma} D^{\lambda\mu j}_{\rho s\sigma}\alpha^{+}_{s\sigma}Q^{+}_{j}(\lambda\mu)\right\}\Psi_0 , \qquad (I)$$

where $Q_j(\lambda\mu)$ is the phonon operator of multipolarity $(\lambda\mu)$, $\alpha^{+}_{s\sigma}$ is the quasiparticle creation operator, $\sigma = \pm 1$, Ψ_0 is the ground state wave function for an even-even nucleus, $\rho\sigma$ stands for the set of quantum numbers specifying the single-particle level with a given K^{π} and $s\sigma$ are the remaining average field levels.

To calculate the energies and the wave functions of the nonrotational states in odd-A nuclei we should find the expectation values of the Hamiltonian involving the pairing and the multiple-multiple interaction. From the condition of the energy minimum we obtain the following secular equation:

$$\varepsilon(\rho) - \eta_i - \frac{1}{4}\sum_{\lambda\mu j}\sum_{s}\frac{v^{2}_{\rho s}}{Y^{j}(\lambda\mu)}\frac{\left[f^{\lambda\mu}(\rho s)\right]^{2}}{\varepsilon(s) + \omega^{\lambda\mu}_{j} - \eta_i} = 0, \qquad (2)$$

defining the energies η_i of the nonrotational states. Here $\varepsilon(s)=\sqrt{C^2+[E(s)-\lambda]^2}$ (C is the correlation function, λ is the chemical potential), $v_{fs}=u_f u_s - v_f v_s$ (u_s and v_s are the coefficients of the Bogolubov transformation), $f^{\lambda\mu}(f s)$ is the matrix element of the multipole moment operator, the quantity $Y^j(\lambda\mu)$ characterizes the collectiveness of the single-phonon state, the explicit form of it is given in refs.[1,2] $\omega_j^{\lambda\mu}$ are the single-phonon state energies. The summation in eq.(2) over $\lambda\mu j$ means that the interaction of quasiparticles with quadrupole $\lambda = 2, \mu = 0,2$ and octupole $\lambda = 3$, $\mu = 0,1,2$ phonons is taken into consideration.

Using the normalization condition for the wave function (I) we get the functions C_f^i and $D_{fs\sigma}^{\lambda\mu i j}$ in the form

$$(C_f^i)^{-2} = 1 + \frac{1}{4}\sum_{\lambda\mu i}\sum_s \frac{v_{fs}^2}{Y^j(\lambda\mu)} \frac{f^{\lambda\mu}(fs)}{\varepsilon(s)+\omega_j^{\lambda\mu}-\eta_i}, \tag{3}$$

$$D_{fs\sigma}^{\lambda\mu i j} = \frac{1}{2}\frac{v_{fs}}{\sqrt{Y^j(\lambda\mu)}} \frac{f^{\lambda\mu}(fs)}{\varepsilon(s)+\omega_j^{\lambda\mu}-\eta_i}. \tag{4}$$

The quantity $(C_f^i)^2$ defines the contribution of the single-quasiparticle component with a given f to the wave function of the state in question. The quantity $(d_{f\sigma s}^{\lambda\mu i j})^2 = \frac{1}{2}(C_f^i)^2\sum_\sigma (D_{fs\sigma}^{\lambda\mu i j})^2$ defines the contribution of the component with a quasiparticle in the s state plus phonon $\lambda\mu i$ to the wave function $\Psi(K^\pi)$.

Each state with given K^{π} has its own version of eq.(2) the solutions of which are the energies η_1 , η_2 , ... For the ground state of an odd-A nucleus $\eta_1(K_0^{\pi_0}) \equiv \eta_F$ takes the lowest value, and the excited state energies are

$$\mathcal{E}_i(F) = \eta_i(K^{\pi}) - \eta_F . \tag{5}$$

The quasiparticle-phonon interactions lead to the appearance of admixtures in the quasiparticle and phonon states and the appearance of complex structure states. The quasiparticle--phonon interactions in odd-A spherical[3] and deformed [1,2,4,5] nuclei are the best investigated ones.

The quasiparticle-phonon interactions lead to small admixtures in the ground and low-lying single-quasiparticle states in odd-mass deformed nuclei. With increasing excitation energy the role of admixtures becomes more important. The quasiparticle-phonon interactions lead to a fragmentation of the single-particle states of the Nilsson and Saxon-Woods potentials in a number of nonrotational levels with given K^{π}. With increasing excitation energy the fragmentation of the single-particle states is performed in an ever-increasing number of levels [6].

The investigation of the structure of odd-A deformed nuclei by means of the (dp) and (dt) reactions showed that the experimentally observed number of the nonrotational states with given K^{π} is much larger than that given by the Saxon-Woods scheme, although the total excitation intensity is in agreement with the estimates obtained with the wave functions of the Saxon-Woods potential. These experiments prove the conclusions of the model about the fragmentation of the single--quasiparticle components in many nonrotational levels.

The degree of fragmentation of the three-quasiparticle components in many levels is expected to be still much larger. Among the three-quasiparticle states, the states with the largest K are pure, the states out of which it is impossible to make the combination quasiparticle plus phonon with $\lambda = 2$ or $\lambda = 3$ are relatively pure.

At excitation energies of 2-3 MeV or higher the fragmentation of the single- and three-quasiparticle components in the nonrotational levels in odd-A nuclei must be essentially increased compared to lower excitation energies. The experimental data on the structure of such states in deformed and spherical nuclei is very poor. One may hope that the construction of mass-separators on particle beams will make it possible to go still further in the study of the levels of medium and heavy nuclei of an energy of 2-5 MeV.

The great progress in neutron spectroscopy made it possible to begin the study of the structure of highly excited nuclear states. By the highly excited states we mean the ones with excitation energies close to the nucleon binding energy or higher. It is interesting to try to treat the highly excited states by means of the semi-microscopic description in the language of the quasiparticle and phonon operators. Such an approach is suggested in ref.[7].

3. The above model which takes into account the quasiparticle-phonon interaction may be useful in understanding how the fragmentation of the single-particle states proceeds and to what extent the level density increases with increasing excitation energy.

The calculations of the energies and the single-quasiparticle components were performed for a number of states of ^{239}U, the neutron binding energy of which B_n is 4.80 MeV. The calculations were performed with the single-particle energies and the wave functions of the Saxon-Woods potential with A = 237 at $\beta_2 = 0.22$ $\beta_4 = 0.08$ [8]. The nonrotational levels of odd--mass nuclei in the actinide region which are calculated in ref. [9] with these single-particle energies and the wave functions are in good agreement with the corresponding experimental data. Since we are interested in the distribution of the single-quasiparticle components over all the excitations up to 5 MeV we should take into account all the poles of the secular equation up to 5 MeV. To this end, we had to take ten roots for each secular equation defining the phonon energies in ^{238}U, i.e. $j = 1,2,3, \ldots 10$.

The results of calculations are given in Tables I and 2. Table I gives the fragmentation of the single-particle states the single-particle energies of which are near the Fermi surface. Table 2 gives the fragmentation of the following single-particle components: $1/2^+[640]$ lying below the Fermi surface energy by 3.5 MeV, $1/2^+[880]$ and $1/2^+[600]$ lying above the Fermi surface energy by 4.4 and 5.2 MeV. In the upper part of the tables are presented six levels which cover mainly the strength of the single-particle components. In the lower part of the tables are given the energies and the single-quasiparticle components for 16 levels near $B_n = 4.80$ MeV .

It is seen from Table I that for the states the single-particle energy of which is close to the Fermi surface energy the first root of the secular equation (2) contains about 90% of these single-particle state. Thus, the low-lying nonrotational states are close to the single-quasiparticle ones. However, even in these cases the single-particle state is distributed over many levels including the levels of an excitation energy of 4-5 MeV. For example, 16 states in the region 4.4 - 5.0 MeV contain (0.1 - 0.6)% of the corresponding single-quasiparticle component.

The complication of the state structure with increasing excitation energy is demonstrated in Table 2. Six levels containing the largest single-particle components contain only 60-70 % of the single-particle state strength. From the comparison of the data given in Table I and 2 it is seen that the sum of the $(C_f^i)^2$ values over 16 levels in the range of 4.4 - 5.0 MeV increases as the single-particle energy approaches the B_n- value.

The growth of the nonrotational level density in ^{239}U with increasing excitation energy is illustrated in Table 3. It is seen from it that the nonrotational level density at an excitation energy of 2.5 MeV or higher increases by a factor of 10–20 compared to the density in the independent quasiparticle model. At an energy of 4.5 MeV the density of the states with $K^{\pi} = 1/2^{-}$ is larger than that for the states with $K^{\pi} = 5/2^{+}$ and $9/2^{-}$ by a factor of (1.5–3).

The experimentally measured average distance between the states with $K^{\pi} = 1/2^{+}$ in odd-mass deformed nuclei in the actinide region, D, is about 10 eV at excitations close to B_n. The calculations performed in the framework of the independent quasiparticle model give the average distance between the levels $D \approx 1$ MeV. Thus, the difference is five orders of magnitude. The calculations by the model taking into account the quasiparticle-phonon interactions for ^{239}U show that for the $K^{\pi} = 1/2^{+}$ states at an excitation energy of 4.5 – 5.0 MeV $D \approx 10$ keV. Thus, the account of the quasiparticle-phonon interaction leads to an increase in the level density near B_n by about a factor of 100 compared to the independent quasiparticle model. However, to obtain agreement with experiment another three orders of magnitude are needed.

To clarify the process of fragmentation of the single-particle states and to study the structure of high-energy states it is necessary to improve the model taking into account the quasiparticle-phonon interactions by adding to the wave function (I) terms like

$$\alpha^{+}_{s\sigma} Q^{+}_{j_1}(\lambda_1 \mu_1) Q^{+}_{j_2}(\lambda_2 \mu_2),$$

$$\alpha^{+}_{s\sigma} Q^{+}_{j_1}(\lambda_1 \mu_1) Q^{+}_{j_2}(\lambda_2 \mu_2) Q^{+}_{j_3}(\lambda_3 \mu_3)$$

etc. and taking into consideration the phonons with $\lambda \geq 4$.

4. To study the structure of the excited states it is useful to introduce a wave function containing a large number of components with different number of quasiparticles. Using such a wave function it is possible to describe the alpha-, beta- and gamma-transition probabilities and the cross section of nuclear reactions.

In ref. [7] a wave function for the highly excited state of an even-even spherical nucleus is constructed. It contains two-, four-, six- and so on quasiparticle components. Among the products of the two operators $\alpha^+_{jm}\,\alpha^+_{j-m}$ being on the same subshell j there are such for which the total angular momentum I is zero. In order that the components $(\alpha^+_{jm}\,\alpha^+_{j-m})_{I=0}$ do not spoil the wave function due to the presence of spurious states, instead them, the pairing vibration phonon operators Ω^+_{ζ} are introduced.

Let us construct, e.g. the wave function for the excited state of an odd N deformed nucleus in the form

$$\Psi_i(K^{\pi}) = \sum_{s\sigma} b^{in}_{K\pi}(s)\,\alpha^+_{s\sigma}\,\Psi_0 + \sum_{\substack{s_1,q_2,q_3\\ \sigma_1,\sigma_2,\sigma_3}}{}' \sum_t b^{in2t}_{K\pi\,\sigma_1\sigma_2\sigma_3}(s_1,q_2,q_3)\,\alpha^+_{s_1\sigma_1}\alpha^+_{q_2\sigma_2}\alpha^+_{q_3\sigma_3}\Psi_0 + \dots + \sum_{s\sigma}\sum_{t,\zeta} b^{in\Omega_\zeta(t)}_{K\pi}(s)\,\alpha^+_{s\sigma}\,\Omega^+_{\zeta}(t)\,\Psi_0 + \dots \qquad (6)$$

The coefficients $b^{i}_{K\pi}$ define the contribution of the corresponding quasiparticle component, i is the number of the excited state with a given K^{π}. The index t = n indicates the neutron and t = p the proton systems. The summation $\sum'_{s_1, s_2, s_3; \sigma_1, \sigma_2, \sigma_3}$ means that the terms $s_1 = s_2$, $s_1 = s_3$, $s_2 = s_3$ are absent and that $E(s_1) < E(s_2) < E(s_3)$. By Ψ_0 we denote the product of the quasiparticle vacua for the neutron and the proton systems, i.e.

$$\Psi_0 = \Psi_0^n \Psi_0^p ,$$

$$\Psi_0^n = \prod_s (u_s + v_s \alpha^+_{s+} \alpha^+_{s-}) \Psi_{00} , \tag{7}$$

where $\alpha^+_{s\sigma}$ is the neutron creation operator, $\alpha_{s\sigma} \Psi_{00} = 0$. The phonon operators determined separately for the neutron and the proton systems are

$$\Omega_i(t) = \frac{1}{2} \sum_q \{ X^i(q) A^+(q,q) - Y^i(q) A(q,q) \} , \tag{8}$$

where $A(q,q) = \sqrt{2}\, \alpha_{q+} \alpha_{q-}$, i is the number of the root of the corresponding secular equation. In just the same way it is possible to construct the wave functions for the excited states of even deformed nuclei.

The wave function (6) is of a very general form. It can be used for the description of both lowlying and highly excited nuclear states. In fact, when only the first component differs from zero the wave function is a single-quasiparticle one. A number of terms in (6) containing the product of two quasiparticle operators can be presented in the form of the phonon operators describing, e.g. in the Tamm-Dancoff approximation, the quadrupole or octupole vibrations.

The wave function (6) can be used for the description of the excited states of an energy of 2-3 MeV or higher up to energies at which the resonances do not yet overlap, i.e. up to excitation energies at which for the states with given K^{π} the condition

$$\Gamma_n^0 \ll D, \tag{9}$$

is valid, i.e. when the total neutron reduced width of the resonance is much smaller than the average distance between the levels D.

To obtain certain integral characteristics the wave function (6) may be useful for the analysis of overlapping resonances and thermal neutron absorption.

If the wave function (6) has a large number of non-zero components, among which there are both few- and many-quasiparticle ones, then it is the wave function of a compound state. The presence of many-quasiparticle components means that the particle penetrated into the nucleus has undergone many collisions by breaking many pais. The small values of the many-quasiparticle components of this wave function result in a hindrance of the probabilities of high-energy gamma transitions to the low-lying states, therefore the half-life of this state must be much longer than that for the single-quasiparticle state.

5. We consider the strength functions for s-neutrons. The cross section for the s-neutron capture and the reduced neutron width for an i-th resonance on an even-even spherical nucleus is written in the form

$$\sigma_i = \sigma_i^{o} \left| b_j^{in}(j) \right|^2 u_j^2 + \sigma_i^{oo}, \tag{10}$$

$$(\Gamma_{no}^{o})_i = \left| b_j^{in}(j) \right|^2 u_j^2 + \Gamma_i^{oo}, \tag{11}$$

$j^{\pi} = 1/2^{+}$, $b_j^{in}(j)$ is the coefficient in the wave function (6). In eq.(10) we singled out the term corresponding to the machanism of the direct neutron capture, σ_i^{o} being the kinematic term. The same is done in the expression for the reduced neutron width. The faotor u_j^2 points out that the contribution to σ_i and $(\Gamma_{no}^{o})_i$ is given by the particle state (i.e. states lying above the Fermi surface).

The reduced width of the s-neutron capture on a deformed even-even nucleus is of the form:

$$(\Gamma_{nc}^{o})_{i} = \left| \sum_{s} b_{1/2^+}^{in}(s)\, u_s \right|^2 + \Gamma_i^{oo}, \qquad (12)$$

The summation over s is connected with the fact that several single-particle $K^{\pi} = 1/2^+$ states may contribute to the wave function (6).

The strength function for the s-neutron is defined as an average value of the neutron width $\langle \Gamma_{nc}^{o} \rangle$ over a number of resonances divided by the average distance D between the levels with $K^{\pi} = 1/2^+$, i.e.

$$S_n = \frac{\langle \Gamma_{nc}^{o} \rangle}{D}. \qquad (13)$$

Using eqs.(11) and (12) we get the strength functions for the s-neutron capture on an even-even spherical and deformed nuclei in the form

$$S_n = \frac{1}{\Delta E} \cdot \sum_{i} \left| b_j^{in}(j) \right|^2 u_j^2 + S_n^{oo}, \qquad (14)$$

$$S_n = \frac{1}{\Delta E} \sum_{i} \left| \sum_{s} b_{1/2^+}^{in}(s)\, u_s \right|^2 + S_n^{oo} \qquad (15)$$

Here the summation i is performed over the resonance number in the energy interval ΔE, $j^{\pi} = 1/2^{+}$. The strength function for the s-neutron capture on an odd N deformed nucleus is

$$S_n = \frac{1}{\Delta E} \sum_{i} \left| \sum_{s_2} b^{i2n}_{K\pi}(s_i, s_2)\, u_{s_2} \right|^2 + S_n^{oo}, \tag{16}$$

where s_i relate to the single-particle level which is occupied by a quasiparticle of the ground state of an odd N target-nucleus. The summation of s_2 is performed over the single-particle states with $K^{\pi} = 1/2^{+}$. The terms S_n^{oo} in eqs.(14), (15), (16) are due to a more complicated (rather than the direct one) mechanism of neutron capture.

According to the widespread opinion (e.g. ref. [10]) in the strength function the term $S_n - S_n^{oo}$ is predominant, i.e. the strength function defines the contribution from simple configurations. Sometimes the neutron strength function is defined as the sum of the spectrosoopic factors of the neutron transfer reaotion of the (dp) type per energy unit.

We give qualitative explanation of the behaviour of the strength function for the s-neutron as a function of the mass number A without recourse to the optical nuclear model. We start from the assumption that the wave functions (6) for the states of an excitation energy close to the neutron binding energy B_n contain a relatively large quasiparticle component with $I^{\pi} = 1/2^+$ when the single-particle neutron level $ns_{1/2}$ in the case of spherical nuclei or $K^{\pi} = 1/2^+$ level in the case of deformed nuclei are near the energy B_n. As the neutron single-particle level with $I^{\pi} = 1/2^+$ moves away from the energy B_n the contribution of the corresponding quasiparticle component to the wave function (6) must decrease.

In the upper part of Fig.I we give the experimental data on the strength functions for s-neutrons $S_n \cdot 10^4$ taken from ref. [11] . In the lowerpart of the same figure we give the behaviour of the neutron states $3s_{1/2}$, $4s_{1/2}$ and [640], 1/2$^+$[651], 1/2$^+$[600], 1/2$^+$[611], 1/2$^+$[880] for deformed nuclei, the energy of these states is reconed from the B_n value (the negative values correspond to the binding states).

Let look for the behaviour of the strength function depending upon A. In the range A = 40 - 60 the state $3s_{1/2}$ is near B_n, therefore the strength function must be maximum. This is in agreement with the available experimental data.

In the range A = 60 - 100 the state $3s_{1/2}$, moves away from B_n remaining above the Fermi surface energy which leads to a decrease of the strength function. In the range A = 100 - 120 the state $3s_{1/2}$ is near the Fermi surface energy and its contribution to the wave functions (6) is expected to be not large. In this range the subshells $3p_{1/2}$ and $3p_{3/2}$ are near B_n and therefore the strength functions must have large values for p-neutrons.

In the range A = 130 - 140 the S_n value increases which is hard to explain since the subshell $4s_{1/2}$ is still above B_n by 6–8 MeV. It is possible that for a number of nuclei with A $\approx$ 100 and A $\approx$ 120–130 being in the transition region some excited states are deformed. For them there are states with $K^{\pi} = 1/2^+$ (marked on Fig.I) close to B_n.

In the range A = 145 – 155 the strength function has a maximum. Then it is assumed that S_n somewhat decreases and there is a second maximum in the range A $\approx$180. The presence of the $1/2^+[640]$ and $1/2^+[651]$ states near B_n makes it possible to account for large values of S_n for deformed nuclei in the region 150 < A < 190. However it is hard to explain simply the decrease of S_n at A $\approx$ 165 compared to the values at A $\approx$ 150. New experimental data (e.g. ref. [12]) point to sharp changes of S_n from isotope to isotope rather than to the minimum near A = 165.

The presence of the subshell $4s_{1/2}$ for spherical nuclei with A = 190 – 220 near B_n points out that the strength functions S_n must have large values.

It is not difficult to explain the large values of S_n in deformed nuclei with A = 226 – 250 since the states $1/2^+[600]$, $1/2^+[611]$ and $1/2^+[880]$ are near the B_n values.

Thus, basing on the location of the single-particle levels with $I^{\pi} = 1/2^+$ one succeeds in explaining qualitatively the behaviour of the strength function for the s-neutron without recourse to the optical nuclear model.

6. Some important information on the structure of highly excited states may be obtained from alpha decay studies. There are experimental data on alpha transitions from the resonanses

to the ground, single-phonon and higher excited states in even-even spherical nuclei [13,14]. In ref. [7] one calculated the reduced widths for the alpha transitions from the highly excited states described by the wave functions (6) to the ground, single-phonon and two-quasiparticle states. It is shown that the alpha decay involves a restricted number of the components of the highly excited state wave function, therefore from the analysis of the appropriate experimental data it is possible to obtain information on the magnitude of these components.

The calculations showed that the alpha decay to the ground state of an even-even nucleus proceeds from the two-quasiparticle, four-quasiparticle of the type 2n2p and two quasiparticles plus phonon Ω_i^+ components of the wave function (6) and all the quasiparticle operators must be particle ones. The reduced probability of the alpha transition from the two-quasiparticle components of the type particle-particle of the wave function (6) is noticeably enhanced. The fluctuation of the particle-particle components in (6) from resonance to resonance must lead to the fluctuation of the corresponding alpha widths.

Table 4 gives the experimental data obtained in ref. [14] on the reduced alpha widths for the transitions from the resonances with $I^{\pi} = 3^-$ in ^{148}Sm to the ground and one-phonon 2^+ states in ^{144}Nd. It is seen from the table that the alpha width ($\gamma_{\alpha 0}^2$)$_i$ for the transitions to the ground state strongly fluctuate from resonance to resonance.

The alpha transitions to the one-phonon states involves a much larger number of the components of the wave function (6) compared to the transitions to the ground states. Therefore the reduced probabilities for the transitions to the one-phonon states must be larger than those to the ground states. This fact is well demonstrated in Table 4.

If the reduced probability for the alpha transition to the ground state is larger than that to the one-phonon one then this indicates that the wave function (6) contains relatively large two-quasiparticle components. Such a feature is characteristic of the resonance with an energy 183.7 eV in ^{148}Sm (Table 4).

The alpha transitions to the two-phonon states involve a much larger number of the components of the wave function (6) compared to the transitions to the one-phonon states. Therefore the reduced probabilities for alpha transitions to them may be larger than those for transitions to the ground and one-phonon states. In ref. [14] one measured the reduced alpha width for the transition from the resonance of an energy of 39.7 eV in ^{148}Sm to the two-phonon 4^+ state in ^{144}Nd, it was found to be $(\gamma^2_{\alpha 2})_i = 1.22$, i.e. by an order of magnitude larger than $(\gamma^2_{\alpha 1})_i = 0.12$ and by two orders of magnitude larger than $(\gamma^2_{\alpha 0})_i = 0.014$. In the same paper it was found that $(\gamma^2_{\alpha 2})_i = 2.44$ and $(\gamma^2_{\alpha 1})_i = 0.35$ for transitions from the resonance with $I^{\pi} = 4^-$ in ^{148}Sm to the two-phonon state 4^+ and the one-phonon 2^+ states. These data indicate to the fact that in some cases there are large alpha widths for the transitions to the two-phonon states.

7. An important information on the structure of highly excited states may be obtained from gamma transitions to the low-lying states. In ref. [7] one calculated the reduced probabilities of the EI and MI transitions from the highly excited states to the ground, one-phonon and two-quasiparticle states in even-even spherical nuclei. It is shown that the gamma transition to the ground state goes from the two-quasiparticle component of the type particle-hole of the wave function (6). The gamma transitions to the one-phonon states involve a large number of components of the wave function (6) compared to the transitions to the ground states.

The gamma transitions from the highly excited states to the two-quasiparticle states involve a large number of the components of the wave function (6). The selection rule requires that one or two quasiparticles in the highly excited state be on the same subshells as the quasiparticles in the final states. Owing to these selection rules it may be expected that the reduced probabilities for gamma transitions to the two-quasiparticle states will be smaller than those to the one-phonon and ground states.

Using the excited state wave function (6) let us analyse which processes may be correlated with one another. The correlation of two processes proceeding throughout the same state takes place if the main contribution to both processes comes from the same components of the wave function of this state

The performed analysis showed that there must be correlations between the following quantities: 1) the reduced neutron widths ($\Gamma_{n i}^{o}$) and the reduced probabilities for the EI transitions to the single-quasiparticle states in odd-mass nuclei and to those two-quasiparticle states of the even spherical nuclei in which one quasiparticle occupies the level corresponding to the ground state of the nucleus which captured the neutron; 2) the reduced neutron widths $(\Gamma_{n i}^{o})_i$ and the reduced EI and MI transitions to the single-quasiparticle and two-quasiparticle (one quasiparticle of which is on the level of the Fermi surface of an odd A-I nucleus) of deformed nuclei.

The examples of such correlations were given in a number of papers, for example, in ref. [12].

There are no correlations between $(\Gamma_{n i}^{o})_i$ and MI transitions in spherical nuclei, $(\Gamma_{n i}^{o})_i$ and EI and MI transitions to the two-quasiparticle proton states and to those two-quasiparticle neutron states not a single quasiparticle of which lie on the level of the Fermi surface of an odd A-I nucleus.

We consider the correlations between the reduced alpha widths $(\gamma_{\alpha n}^{2})_i$ and the reduced neutron widths $(\Gamma_{n i}^{o})_i$ as well as between $(\gamma_{\alpha n}^{2})_i$ and the gamma reduced widths from the same resonance i. The components $b_J^{i 2n}(j_1, j_2)$ of the wave function of the highly excited state of an even-even nucleus enter the expressions for $(\Gamma_{n i}^{o})_i$ and $(\gamma_{\alpha 0}^{2})_i$ which indicates the existence of the correlation between these processes. However, the alpha decay to the ground state involves also a number of four-quasiparticle components and a far larger number of the two-quasiparticle components than in the neutron capture. Therefore the correlation between $(\gamma^2_{\alpha 0})_i$ and $(\Gamma^0_{n0})_i$ can take place only in some cases.

The reduced probabilities of alpha- and EI transitions from the resonance to the ground states of even-even nuclei contain the quantities $b^{2n}_J(j_0, j_2)$ where j_0 related to the subshell which is occupied by a quasiparticle in the ground state of an odd N nucleus A-I. The quantity $b^{2n}_J(j_0, j_2)$ together with the factors $u_{j_0} u_{j_2}$ enter the matrix element of the alpha-decay and with the factor $(u_{j_0} v_{j_2} + v_{j_0} u_{j_2})$ in the matrix element of the EI transition. In the case $j = j_0$ it may be approximately assumed that $u_{j_0} = v_{j_0}$. Therefore the existence of the correlation between both processes is possible. However, the expression $(\gamma^2_{\alpha 0})_i$ contains many other components of the wave function of the highly excited state.

Of great interest is the experimental determination of the quantities $(\gamma^2_{\alpha n})_i$ and $(\Gamma^0_{nc})_i$ and the reduced probabilities for gamma transitions for the same resonances. It would be desirable to determine $(\Gamma^0_{n0})_i$ and the probabilities of gamma transitions for the resonance $I^{\pi} = 3^-$ with an energy 183.7 eV in ^{148}Sm which has a rather large width $(\gamma^2_{\alpha 0})_i$.

The performed investigations showed that the expression of the wave function in the form (6) turns out to be useful for the study of the structure of these states.

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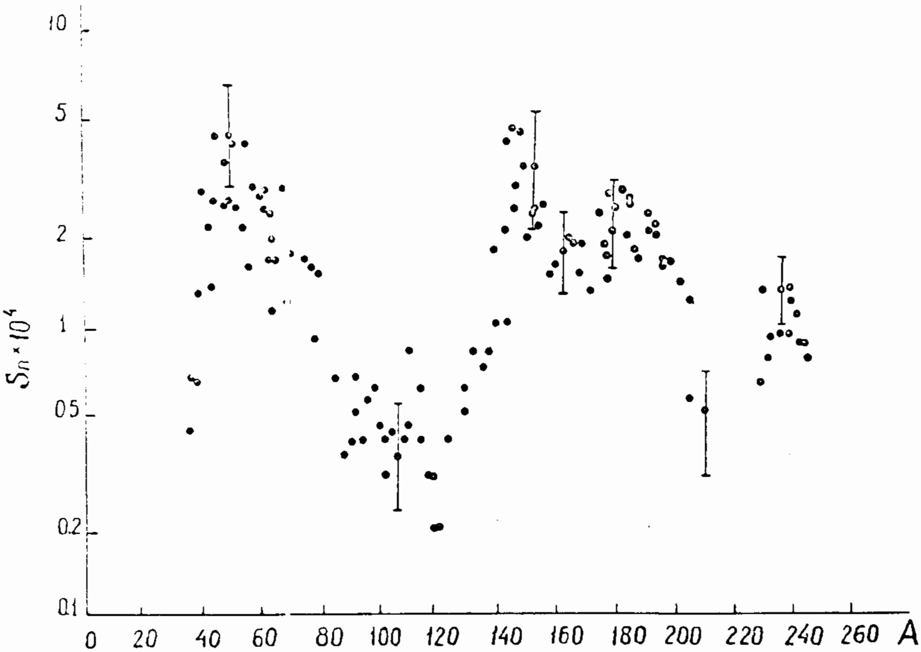
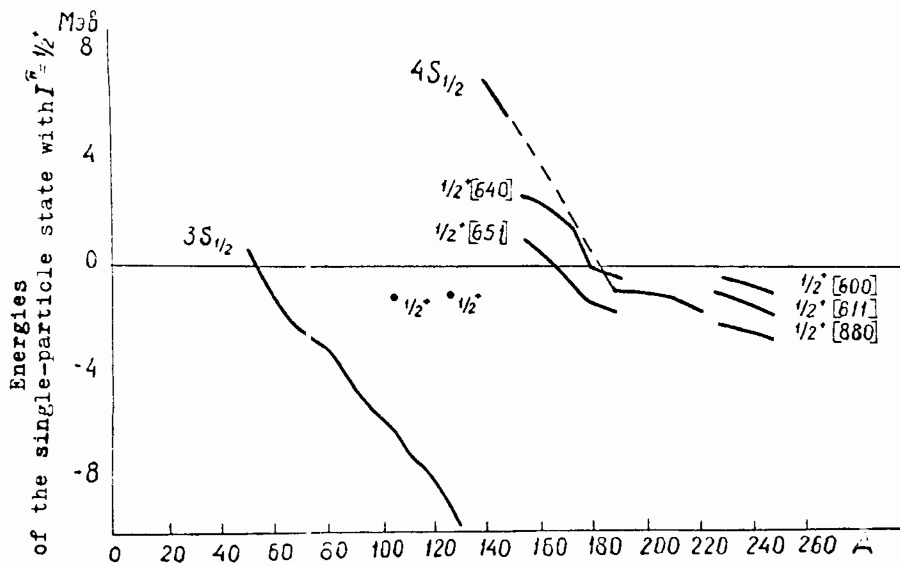


10
5
2
1
0.5
0.2
0.1
$S_n \times 10^4$
0 20 40 60 80 100 120 140 160 180 200 220 240 260 A



Energies
of the single-particle state with $I^{\pi}=1/2^{+}$
Мэв
8
4
0
-4
-8
$4S_{1/2}$
$3S_{1/2}$
1/2⁺[640]
1/2⁺[651]
1/2⁺ 1/2⁺
1/2⁺[600]
1/2⁺[611]
1/2⁺[880]
0 20 40 60 80 100 120 140 160 180 200 220 240 260 A

Fig. 1

T A B L E I

Fragmentation in ^{239}U of the single-particle states near the Fermi surface energy.

5/2−[622], $\varepsilon(s_i)$=0,712		1/2−[631], $\varepsilon(s_i)$=0,812		9/2−[734], $\varepsilon(s_i)$=1.076	
$\varepsilon_i - \varepsilon_F$ MeV	$(C^{i}_{s_i})^2\%$	$\varepsilon_i - \varepsilon_F$ MeV	$(C^{i}_{s_i})^2\%$	$\varepsilon_i - \varepsilon_F$ MeV	$(C^{i}_{s_i})^2\%$
0	91.7	0.102	88.7	0.245	85.5
1.784	0.6	1.260	0.8	1.559	3.3
1.884	1.4	1.983	1.5	1.639	3.0
2.892	0.4	2.504	1.4	2.249	1.2
2.906	1.4	2.768	1.1	2.298	1.6
3.371	0.5	4.700	0.6	5.992	2.0
Sum	96.0	Sum	94.1	Sum	96.6
4.690	$1\cdot 10^{-5}$	4.768	$3\cdot 10^{-5}$	4.401	0.00117
4.722	$1\cdot 10^{-6}$	4.769	0.14360	4.426	0.00706
4.737	$7\cdot 10^{-5}$	4.792	0.40573	4.432	0.00014
4.746	$1\cdot 10^{-6}$	4.801	0.06074	4.450	0.00026
4.765	$2\cdot 10^{-5}$	4.807	$< 10^{-6}$	4.471	0.00089
4.770	$2\cdot 10^{-6}$	4.808	$1\cdot 10^{-6}$	4.520	$9\cdot 10^{-5}$
4.825	$1\cdot 10^{-5}$	4.818	$6\cdot 10^{-6}$	4.533	$2\cdot 10^{-5}$
4.833	$4\cdot 10^{-6}$	4.826	0.00065	4.541	$1\cdot 10^{-5}$
4.892	$< 10^{-6}$	4.857	$4\cdot 10^{-5}$	4.585	$9\cdot 10^{-5}$
4.898	$5\cdot 10^{-5}$	4.862	0.00074	4.676	0.00044
4.910	$4\cdot 10^{-5}$	4.871	0.00011	4.689	0.00971
4.912	0.00821	4.873	$1\cdot 10^{-5}$	4.[illegible]25	$6\cdot 10^{-5}$
4.917	0.08576	4.901	$5\cdot 10^{-5}$	4.[illegible]67	0,07692
4.922	0.42494	4.916	$8\cdot 10^{-5}$	4.[illegible]70	0.00168
4.930	0.00024	4.934	$< 10^{-6}$	4.[illegible]84	0.00120
4.931	0.00576	4.935	$4\cdot 10^{-5}$	4.[illegible]73	0.00018
Sum	0.52512	Sum	0.61181	Sum	0.0999[illegible]

T A B L E 2

Fragmentation in ^{239}U of the single-particle states far away of the Fermi surface energy.

1/2 +[640], $\varepsilon(s_i)$= 3,530		1/2 +[880], $\varepsilon(s_i)$= 4,424		1/2+[600], $\varepsilon(s_i)$=5,251	
$\eta_i - \eta_F$, MeV	$(C^i_{s_i})^2\%$	$\eta_i - \eta_F$, MeV	$(C^i_{s_i})^2\%$	$\eta_i - \eta_F$, MeV	$(C^i_{s_i})^2\%$
1.332	2.7	2.585	16.0	2.650	4.6
1.890	34.4	2.587	5.4	2.937	35.4
2.351	13.8	2.827	11.0	3.058	12,1
2.890	5.4	2.861	14.5	3,097	2,5
4.081	10,6	3.108	14,4	3,285	0.2
4.508	3.8	4.560	1.6	4.680	4.1
Sum	70.7	Sum	62,9	Sum	58.9
4.768	0.00203	4.768	10^{-6}	4,768	$2\cdot10^{-5}$
4.769	1.11736	4.769	10^{-6}	4.769	0,00031
4.792	0.09847	4.791	$2\cdot10^{-6}$	4.792	$1\cdot10^{-6}$
4.801	0.02320	4.801	$1\cdot10^{-6}$	4,802	0,86308
4.807	$2\cdot10^{-5}$	4.806	$4\cdot10^{-5}$	4.804	0.00012
4.808	0.00079	4.808	$<10^{-6}$	4.810	0.01214
4.818	$6\cdot10^{-6}$	4.819	0.38851	4.818	$1\cdot10^{-6}$
4.826	$18\cdot10^{-5}$	4.828	0.25372	4.831	0.09227
4.857	$29\cdot10^{-5}$	4.857	$4\cdot10^{-6}$	4.857	$2\cdot10^{-5}$
4,862	$4\cdot10^{-6}$	4.861	0.02102	4.860	0.27764
4.871	$7\cdot10^{-5}$	4.871	$1\cdot10^{-5}$	4.871	0.00080
4,873	0.00051	4.873	$1\cdot10^{-6}$	4.873	0.00052
4.901	$3\cdot10^{-5}$	4.900	0.00211	4.894	3.25660
4.916	0.00639	4.916	$5\cdot10^{-5}$	4,916	$8\cdot10^{-5}$
4.934	$<10^{-6}$	4.932	0.08336	4.934	$<10^{-6}$
4,934	0.00352	4.935	$6\cdot10^{-5}$	4.935	0,00010
Sum	1.252	Sum	0.749	Sum	4.502

T A B L E 3

Number of non-rotational states in ^{239}U

Excitation energy interval MeV	Number of nonrotational states with		
	$K^{\pi} = 1/2^{+}$	$K^{\pi} = 7/2^{+}$	$K^{\pi} = 9/2^{-}$
0 - 0.5	1	1	1
0.5 - 1.0	0	0	0
1.0 - 1.5	0	3	1
1.5 - 2.0	2	7	2
2.0 - 2.5	5	11	11
2.5 - 3.0	13	29	22
3.0 - 3.5	18	17	40
3.5 - 4.0	27	23	17
4.0 - 4.5	42	31	14
4.5 - 5.0	50	30	11

T A B L E 4

Reduced widths for alpha transitions from the states with $I^{\pi} = 3^-$ in ^{148}Sm to the ground and one-phonon states of ^{144}Nd

Energy of resonance $\varepsilon - B_n$, eV	Reduced widths for transition to states ground $(\gamma^2_{\alpha 0})_i$	one-phonon $(\gamma^2_{\alpha 1})_i$
3.4	0.425	1.14
29.7	0.098	0.11
39.7	0.014	0.12
83.3	0.480	0.86
102.6	0.044	0.99
123.4	0.191	0.25
183.7	5.720	-