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THE STRUCTURE OF NUCLEAR STATES AT LOW, INTERMEDIATE AND HIGH EXCITATION ENERGIES

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Summary

It is shown that within the model based on the quasiparticle-phonon interaction one can obtain the description of few-quasiparticle components of nuclear states at low, intermediate, and high excitation energies. For the low-lying states the energy of each level is calculated. The few-quasiparticle components at intermediate and high excitation energies are represented to be averaged in certain energy intervals and their characteristics are given as the corresponding strength functions. The fragmentation of single-particle states in deformed nuclei is studied. It is shown that in the distribution of the single-particle strength alongside with a large maximum there appear local maxima and the distribution itself has a long tail. The dependence of neutron strength functions on the excitation energy is investigated for the transfer reactions of the type (d,p) and (d,t). The s,- p,- and d- wave neutron strength functions are calculated at the neutron binding energy Bn. A satisfactory agreement with experiment is obtained. The energies and $E\lambda$ -strength functions for giant multipole resonances in deformed nuclei are calculated. The energies of giant quadrupole and octupole resonances are calculated. Their widths and fine structure are being studied. It is stated that to study the structure of highly excited states it is necessary to find the values of many-quasiparticle components of the wave functions. The ways of experimental determination of these components based on the study of γ -transitions between highly excited states are discussed.

1. The basic statements of the semi-microscopic description of the nuclear structure are the following:

1) The Hartree-Fock-Bogolubov method is used to obtain the closed system of equations for the density and correlation functions. This is the basic approximation in nuclear many-body problem.

2) The representation in which the density matrix is diagonal and the correlation function is of the canonical form is chosen.

3) The average field described by the Saxon-Woods potential corresponds to the above representation for a number of dubly even nuclei in the beta-stability zone. All residual interactions in such a case are reduced to pairing superconducting interactions. The average field defines many nuclear properties directly and governs the residual interactions.

4) The excitation states are defined as one-, two-, three- and many- quasiparticle states.

5) The low-lying vibrational states are connected with non-diagonal parts of the density matrix. To describe them the multipole-multipole and spin-multipole-spin-multipole interactions are introduced. The mathematical treatment is based on various versions of the approximate second quantization method, first developed by N.N.Bogolubov /1/.

c) The rotational, quasiparticle and phonon excitation states are coupled between themselves by the Coriolis and quasipartical phonon interactions.

When describing low-lying non-rotational states, the nuclear Hamiltonian is written as:

$$H = H_{av} + H_{pair} + H_{Q} \tag{1}$$

It is added by the spin-multipole-spin-multipole interactions, the Gamow-Teller type interactions, and a number of others. The study of the low-lying states allowed one to fix the parameters of the Saxon-Woods potential, pairing constants, isoscalar constants of quadrupole-quadrupole $\varkappa_0^{(2)}$ and octupole-octupole $\varkappa_0^{(3)}$ interactions/2-4/. A sufficiently good description of the low-lying excitation states has been obtained within the semimicroscopic method.

In this report we shall analyze whether the nuclear interactions, given as (1) and the developed methods of solution of nuclear many-body problem can serve as the base for the description of the structures of low, intermediate and high excitation states.

2. A large region of intermediate and high excitation energies of an atomic nucleus lies between the low-lying states and the states which may be described by the extreme statistical model. The level density here is large, and the state structure is highly complicated.

The experimental study of the state structure of this region encounters great difficulties. It is practically impossible to measure the characteristics of each of many thousands levels. Moreover, with increasing excitation energy the state structure becomes complicated, and therefore, there increases the number of components of the corresponding wave functions which should be measured experimentally. However, one should explain the main properties of these states. In the region of intermediate excitation energies, new branches of the collective excitations, for example, may be detected. It is a common viewpoint that the strength of possible new collective branches will not be concentrated mainly on one level but distributed in a certain energy interval. Possibly, there are collective excitation branches corresponding to high multipoles and spin-multipoles. If there were, for instance, strongly collectivized spin-multipole states with $\lambda \geqslant 5$ in the range of intermediate excitation energies in heavy nuclei, we would hardly detect them up to now. It should be noted that the whole experimental information is limited to few quasiparticle components of the corresponding wave functions /5/ unless we obtain the quantities like the spectroscopic factors of the multinucleon transfer reactions and measure γ -transitions between the high states and so on.

The complication of the state structure proceeds already at low excitation energies, this can be illustrated by many spherical and transition nuclei. In ref. /6/ it is shown that in odd-A deformed nuclei at excitation energies 1 MeV and higher, an important role is played by the components quasiparticle plus two phonons of the wave function. The method based on the quasiparticle-phonon interaction with the wave function in the form:

$$\Psi_i = \left\{ \sum_t C_t^i Q_t^+ + \sum_{t_1 t_2} D_{t_1 t_2}^i Q_{t_1}^+ Q_{t_2}^+ \right\} \Psi_0 \qquad (2)$$

was used in ref. /7/ to calculate anharmonic effects. Here Q_i is the phonon operator, $i = (\lambda\mu j)$ j the root number of the phonon secular equation, Ψ_0 is the wave function of the ground state. Let us use this model and calculate the energies and wave functions of five lower excitation states with $K^{\pi} = 0^+, 2^+, 0^-, 1^-$ and 2^- in a number of deformed nuclei and then trace the change of their structures with slight change of $\varkappa_0^{(2)}$ and $\varkappa_0^{(3)}$. Table I shows that with increasing $\varkappa_0^{(3)}$ by 2.7%, the structure of all states but the first ones changes strongly. The investigations showed that the structure of the first and a number of the second states with $K^{\pi} = 0^+, 2^+, 0^-, 1^-, 2^-, 3^-$ changes little, whereas the structure of most of the third and higher states changes very strongly.

Table I

Change of the Energy and Structure of the First Five $K^{\pi} = 1^-$ States in ^{158}Gd With Increasing $\varkappa_0^{(3)}$ by 2.7%

	$\varkappa_0^{(3)} = 0{,}4646 \ 10^{-4} MeV/(fm)^6$			$\varkappa_0^{(3)} = 1{,}027 \varkappa_0^{(3)}$		
	$\mathcal{E}$ MeV	Structure, %		$\mathcal{E}$ MeV	Structure, %	
1.	1.146	(311)98.06	(201,311)1.0	0.904	(311)95.0	(201,311) 3.3
2.	2.248	(312)94.2	(201,311)2.9	2.122	(312)20.7	(201,311)31.2
		(313)1.4			(313)3.8	(221,311)43.6
3.	2.434	(313)85.9	(201,311)4.3	2.212	(312)12.1	(201,311)31.3
4.		(312)3.2	(221,311)5.0		(311)3.5	(221,311)53.0
4.	2.579	(313)4.7	(221,311)93.8	2.295	(312)65.9	(201,311)31.8
5.	2.605	(313)5.7	(201,311)91.1	2.442	(313)91.5	(201,311)2.2
		(312)1.9			(201,313)2.7	(211,311)1.9

Thus, one may conclude that it is impossible to describe correctly the structure of each nuclear level at the excitation energy higher than 2-3 MeV within the framework of the existing theories.

To study the state structure of intermediate and high excitation energies and to describe them by the quasiparticle and phonon operators the important role is attributed to the fragmentation, i.e., the distribution of the strength of single-particle, two-particle and many-particle states over many nuclear levels.

To study the state structure on intermediate and high excitation energies, one should clarify the general regularities of the fragmentation of one-, two- and many-quasiparticle states. The one-nucleon transfer reactions are the important tool in the study of the fragmentation of one-quasiparticle states at intermediate excitation energies. First of all it is necessary to measure experimentally and to describe theoretically the strength functions of the one-nucleon transfer reactions which provide information on one-quasiparticle components averaged over several excited states. Due to the coupling of different channels, one cannot extract correctly the spectroscopic factors from experimental data. Thus, one should perform cumbersome calculations of cross-sections. In ref.[8] an attempt was made to obtain information on neutron strength functions in deformed nuclei. Of most interest is the experimental measurement of strength functions for the one-nucleon transfer reactions with the fixed transfer angular momentum ℓ or to the final states with the fixed I . One does not need high energy resolution for such experiments.

3. Our investigation of the state structure of intermediate and high excitation energies is performed in two directions: general consideration based on the operator form of the wave function and calculation of the state characteristics within the model based on the quasiparticle-phonon interaction. The model is described in ref. /9/, in ref. /10/ the approximate methods of solving its equations were developed. In ref. /11/ the model was generalized to the case of spin-multipole forces. In ref. /12/ it is applied for the description of doubly even deformed nuclei, while in ref. /13/ for odd-A spherical nuclei.

The wave function of the model corresponding to an odd-A deformed nucleus is the following:

$$\Psi_i(K^\pi) = \frac{1}{\sqrt{2}} \sum_\sigma \Big\{ \sum_\rho C_\rho^i \alpha^+_{\rho\sigma} + \sum_g D_g^i (\alpha^+ Q^+)_g + \frac{1}{\sqrt{2}} \sum_G F_G^i (\alpha^+ Q^+ Q^+)_G \Big\} \Psi_0 , \qquad (3)$$

where i is the number of the state, $g = \nu t$, $G = \nu t, t_2, (\rho\sigma), (\nu\sigma)$ - denote the quantum numbers of single-particle states. Using the variational principle we have found /10/ the system of main equations and also for $\rho \neq \rho_c$ (ρ_c is a selected single-particle state)

$$\frac{C_\rho^i}{C_{\rho_c}^i} = \frac{\theta_{\rho_c}(\rho; \eta_i)}{\theta_{\rho_c}(\eta_i)} , \qquad (4)$$

where the denominator of (4) is the determinant of the system and the numerator is the determinant in which ρ -th column is replaced by that of free terms. The quantity $(C_{\rho_0}^i)^2$ is determined from the wave function normalization (3). The secular equation for defining the energies η_i symbolically can be written as

$$\mathcal{F}_{\rho}(\eta_i) = 0. \qquad (5)$$

To describe highly excited states within the framework of the model the phonons of multipole and spin-multipole type with $\lambda = 1, \ldots, 7$ and higher as well as a large number of phonons of each multipolarity are taken into account. Alongside with the known low-lying collective quadrupole and octupole phonons we consider many weakly collectivized phonons as well as high-lying phonons like the giant resonances. The configurational space of the model is large, the wave functions of highly excited states comprise millions of different components. Good description (taking into account many-phonon components) of the density of highly excited states proves the completeness of the configurational space /14/.

Some results on the fragmentation of single-particle states in odd-A deformed nuclei were obtained in ref. /15,16/ in the framework of the simplified model with $F = 0$ in eq.(3). Such calculations are represented in fig.1 by a histogram. The sum of quantities over i lying in the energy interval ΔE = 0.4 MeV is denoted by $C_{\rho}^2 = \sum_{i}^{\Delta E} (C_{\rho}^i)^2$ and given in per cent.

To study the fragmentation of single-particle states, one should calculate the energies and wave functions of many states, and sum up $(C_{\rho}^i)^2$ in some energy interval. Thus, only a small part of the obtained results is used. In ref. /16/ the direct calculation method of averaged characteristics without a detailed calculation of each state was used. We construct the function

$$\Phi_{\rho}(\eta) = \sum_{i} (C_{\rho}^i)^2 \rho(\eta_i - \eta), \qquad (6)$$

where

$$\rho(\eta_i - \eta) = \frac{1}{2\pi}\,\frac{\Delta}{(\eta - \eta_i)^2 + (\Delta/2)^2}\ .$$

The energy interval of averaging Δ is a free parameter. We write the function $\Phi_\rho(\eta)$ as:

$$\Phi_\rho(\eta) = -\sum_i \left(\frac{\partial \mathcal{F}_\rho(\eta)}{\partial \eta}\right)^{-1}_{\eta=\eta_i} \rho(\eta_i - \eta)\ . \qquad (6')$$

The function (6') may be written as a contour integral around the poles which are the roots of eq.(5). Considering that the contour integral over infinite radius circle in the complex plane z is equal to zero, we pass to two contour integrals around the poles $z_1 = \eta + i\,\Delta/2$ and $z_2 = \eta - i\,\Delta/2$. We calculate the corresponding residues and get:

$$\Phi_\rho(\eta) = \frac{1}{\pi}\,\mathrm{Im}\left(\frac{1}{\mathcal{F}_\rho(\eta + i\Delta/2)}\right) \qquad (7)$$

For a simplified model with $F=0$ in eq.(3), we have:

$$\Phi_\rho(\eta) = \frac{\Delta}{2\pi}\,\frac{\Gamma(\eta)}{(\varepsilon_\rho + \gamma(\eta) - \eta)^2 + (\Delta/2)^2\Gamma^2(\eta)} \qquad (8)$$

The Breit-Wigner form is strongly distorted due to the dependence on η $\Gamma(\eta)$ and $\gamma(\eta)$ which explicit form is given in ref.[16]. It results in an essential difference from the accepted description of neutron strength functions[17,18]. A more general consideration of strength functions was performed in ref.[19].

The function $\Phi_\rho(\eta)$ calculated by formula (8) with Δ = 0.4 MeV for the 501↓ state in ^{239}U is represented in fig.1 as a curve. It is seen from the figure that due to the dependence of $\Gamma(\eta)$ and $\gamma(\eta)$ on η, the general Breit-Wigner form $\Phi_\rho(\eta)$ is

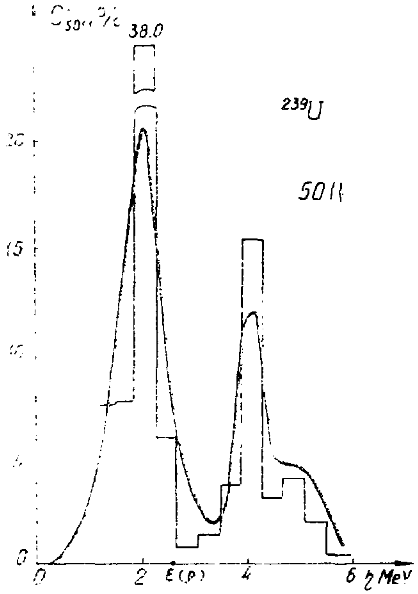

Fig.1. Comparison of C_ρ^i (histogram) and $\Phi_\rho(\eta)$ (curve) for the state 501↑ in ^{239}U calculated with Δ =0.4 MeV. $E(\rho)$ is the quasiparticle energy of the state ρ =501↑ reckoned from the ground state energy (the Fermi 622↑ level).

strougly distorted. The function $\Phi_\rho(\eta)$ follows the behaviour of C_ρ^2 . A similar picture is observed in other cases. Comparing the behaviour of $\Phi_\rho(\eta)$ and C_ρ^2 , it may be concluded that for calculation of the fragmentation of single-particle states in odd-A deformed nuclei the corresponding calculations may be performed by eqs. (7) and (8). Fig.2 shows the fragmentation in ^{171}Er of the $510\uparrow$ state, lying near the Fermi-level, the hole state $402\downarrow$ and deep hole state $404\downarrow$. They are calculated taking account of other single-particle states with given K^π , lying in the energy interval from 8 MeV below and 8 MeV higher than the Fermi-level. The examples of the fragmentation of single-particle states for deformed nuclei are given in refs.[6,15,16,20] and for spherical nuclei in ref.[21].

Thus, the fragmentation of single-particle states in deformed nuclei displays the following features:

i) If the single-particle state is near the Fermi-level, then 90% of the strength is concentrated on the lowest level with the given K^π and the remaining 10% are distributed in a large energy interval;

ii) As the single-particle level moves away from the Fermi-level, the strength concentrated on one level decreases, and the distribution itself expands;

iii) At high quasiparticles energies, in addition to the first large maximum there appear other maxima;

iv) The distribution function is non-symmetric with respect to its largest value due to its slower fall in favour of high energies;

v) the shape of the distribution function is mainly defined by the position of the state with respect to the Fermi-level, it

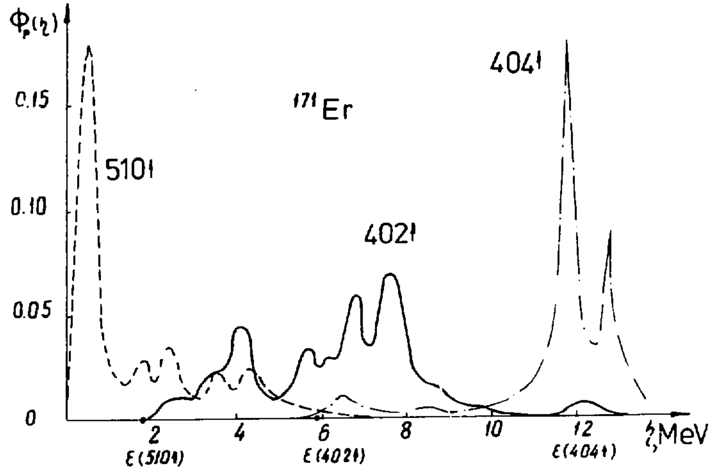

Fig.2. Fragmentation of single-particle states 510↓ (dashed curve) 402↓ (continuous curve) and 404↓ (dash-dotted curve) in ^{171}Er .

also depends on the wave function of the single-particle state and on the influence of other states with the same value of K^{π};

vi) The strength distribution has a long tail which even for the single-particle states lying near the Fermi-level expands father than the neutron binding energy.

Thus, the strength distribution of the single-particle state in a deformed nucleus has a complex dependence on the excitation energy in comparison with that accepted intuitively (refs. /17,18/).

4. The possibility of calculating of the fragmentation of single-particle states has resulted in a new semi-microscopic method of calculation the strength functions for the neutron resonances and those for one-neutron or proton transfer reactions.

The strength function for the one-neutron transfer reaction like (dp) with the fixed ℓ on the doubly even target can be written as:

$$S_{\ell}^{\rho} = S_{\ell,\ell+1/2}^{\rho} + S_{\ell,\ell-1/2}^{\rho} . \tag{9}$$

We neglect the Corioles forces and other interactions resulting in the coupling of various rotational bands what can be achieved provided they do not redistribute essentially the strength of single-particle states from one energy interval ΔE to another. Then

$$S_{\ell,\ell\pm1/2}^{\rho} = \sum_{K} \frac{1}{\Delta E} \sum_{i \in \Delta E} \Big| \sum_{\rho} a_{\ell,\ell\pm1/2}^{\rho K} U_{\rho} C_{\rho}^{i} \Big|^2 = \sum_{K} S_{\ell,\ell\pm1/2}^{\rho K} , \tag{10}$$

where

$$S_{\ell I}^{\rho K}(\eta) = \frac{1}{\pi} \sum_{\rho} (a_{\ell I}^{\rho K} U_{\rho})^2 \, Im \left\{ \frac{1}{\mathcal{F}_{\rho}(\eta + i\Delta/2)} \right\} + \tag{11}$$

$$+ \frac{2}{\pi} \sum_{\rho > \rho'} a_{\ell I}^{\rho K} a_{\ell I}^{\rho' K} U_{\rho} U_{\rho'} \, Im \left\{ \frac{\theta_{\rho}(\rho', \eta + i\Delta/2)}{\theta_{\rho}(\eta + i\Delta/2)} \right\} .$$

Here $a_{\ell I}^{\rho\lambda}$ are the expansion coeffecients of the single-particle wave function of a deformed nucleus in spherical basis[22]. The strength function S_ℓ^h of the reaction like (dt) can be obtained if one replaces u_ρ by v_ρ^2 in eqs. (10) and (11) where u_ρ , v_ρ are the Bogolubov transformation coeffecients.

The calculations of the fragmentation of single-particle states and neutron strength functions showed[16] that besides strongly collectivized phonons one should consider a large number of phonons of each multipolarity from λ =2 to λ =7 since the phonon space can not be strongly limited due to a large number of weakly collectivized phonons.

We calculate the strength functions S_ℓ^p and S_ℓ^h for many nuclei[16,23] including those studied in ref.[8]. Fig.3 represents the functions S_ℓ^p for ℓ =0,1,2 calculated for ^{177}Yb with Δ =0.4 MeV. It is difficult to compare directly our results with experimental data and calculations of ref. [8], since in ref.[8] the cross-sections of (dp) reactions but not the strength functions are considered. However, our results agree in general with those of ref.[8]. Comparing ^{177}Yb with a similar one in ref.[8], one can see that the position of maxima in both cases practically coincide and we describe qualitatively correctly the shape of the experimental curve. Our results on the fragmentation of single-particle states can be used to calculate the cross-sections for (dp) and (dt) reactions, taking into account coupling of different channels.

The fragmentation of single-particle states in spherical nuclei is investigated experimentally by studying the reactions as (dp) [24], $(d\,{}^3He)$, $({}^3He, \alpha)$ [25] and others. The fragmentation

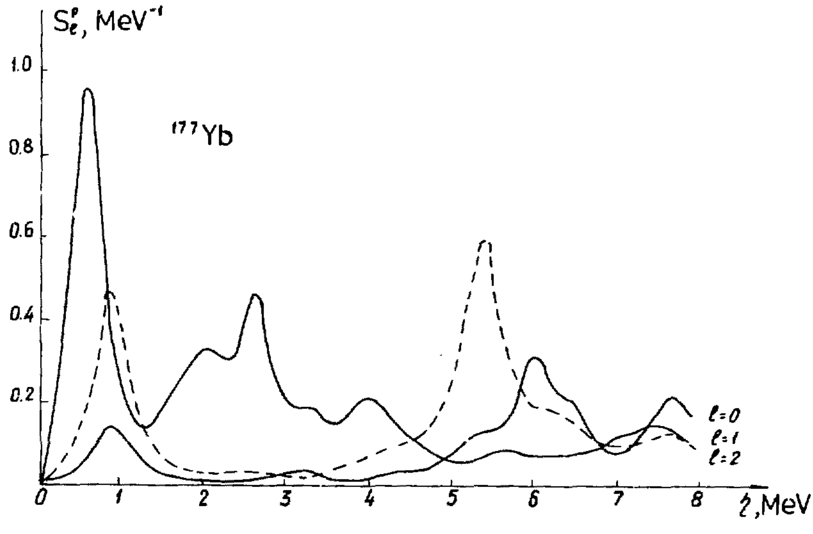

Fig.3. Function S_ℓ^p with $\ell=0$ (thick curve), $\ell=1$ (thin curve) and $\ell=2$ (dash-dotten curve) for ^{177}Yb, calculated with Δ =0.4 MeV.

of subshells in spherical nuclei and the neutron strength functions are calculated within the framework of our model/21/.

The study of the fragmentation of three-quasiparticle states in spherical nuclei is of most interest. There are first experimental data on the three-nucleon transfer reactions in complex nuclei/26/. We hope that our model will serve as a base for the calculation of the fragmentation of three-quasiparticle states in spherical and deformed nuclei.

5. Our mathematical apparatus can be applied for the calculation of neutron strength functions at the neutron binding energy and for comparison of them with the corresponding experimental data. In our notations the s -, p -, and d - wave neutron strength functions, defined in refs./17,18,27/, are the following:

$$S_0 = (15\,KeV)A^{-1/3}S^{\rho 1/2}_{0\,1/2}\ , \tag{12}$$

$$S_1 = (5\,KeV)A^{-1/3}\left\{S^{\rho 1/2}_{1,1/2} + S^{\rho 1/2}_{1,3/2} + S^{\rho 3/2}_{1,3/2}\right\}, \tag{13}$$

$$S_2' = (9\,KeV)A^{-1/3}\left\{S'^{\rho 1/2}_{2,3/2} + S^{\rho 1/2}_{2,5/2} + S'^{\rho 3/2}_{2,3/2} + S^{\rho 3/2}_{2,5/2} + S^{\rho 5/2}_{2,5/2}\right\}, \tag{14}$$

where $S^{\rho K}_{\ell,\ell\pm 1/2}$ is defined by eq(11) and given in units $(keV)^{-1}$. For spherical nuclei S_0 is defined by eq.(12) and

$$S_1 = (5\,KeV)A^{-1/3}\left\{S^{\rho}_{1,1/2} + 2S^{\rho}_{1,3/2}\right\}, \tag{13'}$$

$$S_2' = (9\,KeV)A^{-1/3}\left\{2S'^{\rho}_{2,3/2} + 3S'^{\rho}_{2\,5/2}\right\}, \tag{14'}$$

The s - and p - wave neutron strength functions were calculated for some tin and tellurium isotopes in ref.[21] A good description of the -wave strength in its minimum has been obtained. In the tellurium isotopes the subshell $3S_{1/2}$ is below the Fermi level and with increasing A some increase of C_{ρ}^{2} is compensated by decrease of U_{1}^{2}. A satisfactory description of the ρ -wave neutron strength function in the region close to its maximum is given (in ref.[2] in eqs.(10),(10') the multiplier 1/2 is omitted). One should note, that when calculating the fragmentation and strength functions in spherical nuclei it is necessary to use the model where the quasiparticle plus two phonons are taken into account in the wave function.

The numerical results for the s -, p -, and d - wave neutron strength functions are given in table 2. The experimental data are taken from ref.[28]. The calculations were performed with the parameter Δ =0.4 MeV, in many cases the results slightly depend on the Δ -value. Table 2 shows that a rather good description of the s -, and p -wave neutron strength functions is obtained. We also present the results of the d -wave strength functions for which there are only very preliminary experimental data. The agreement of our results with experimental data is not trivial as the calculations are based on the fragmentation of single-particle states and have no free parameter. Note, that the accuracy of our calculations is limited not only by the accuracy of the approximate description of the fragmentation of single-particle states but also by that of the calculation of single-particle energies and wave functions of the Saxon-Woods potential.

Table 2.

Neutron strength functions at [illegible] B_n

Compound nucleus	B_n MeV	$S_0 \cdot 10^4$ Exp.	Calc.	$S_1 \cdot 10^4$ Exp.	Calc.	$S_2 \cdot 10^4$ Calc.
^{155}Sm	5.819	1.8±0.5	1.0		1.1	1.2
^{159}Gd	6.031	1.5±0.2	1.0	$2.8^{+1.4}_{-1.0}$	1.6	1.0
^{161}Gd	5.650	1.8±0.4	0.9	$0.88^{+0.84}_{-0.47}$	1.1	1.2
^{163}Dy	6.253	1.88	1.8	1.4	0.7	3.7
^{165}Dy	5.635	1.7	1.8	1.3	0.6	3.6
^{169}Er	5.997	1.5	4.0	0.7	0.5	6.8
^{171}Er	5.676	1.54	3.5	0.8	0.7	5.2
^{183}W	6.187	2.1±0.3	4.6	0.3±0.1	0.8	2.0
^{231}Th	5.09	1.3	1.1	-	0.7	4.0
^{233}Th	4.96	0.9	0.8	0.5–1.6	0.5	6.0
^{233}U	5.88	0.95	0.9	-	0.8	4.0
^{235}U	5.27	1.13±0.4	1.3	-	1.2	5.8
^{237}U	5.30	1.3±0.2	1.2	2.3±0.6	1.1	4.6
^{239}U	4.78	1.1±0.1	1.5	1.7±0.3	0.6	3.8
^{241}Pu	5.41	0.94±0.09	0.9	2.8	1.0	3.4
^{243}Pu	5.05	0.9±0.1	1.4	-	1.4	4.0
^{245}Cm	5.696	1.1±0.2	1.6	-	0.7	3.0

6. Recently great progress has been achieved in the study of giant multiplole resonances /29-31/. Theoretical investigations

beyond the framework of the phenomenological method became widely used now [32-34]. Let us give the calculation results for the energies and strength functions of giant multipole resonances performed with the framework of the above model.

The secular equations for the description of one-phonon states with the multipole-multipole interaction constants $\varkappa_{nn}^{(\lambda)} = \varkappa_{pp}^{(\lambda)} \neq \varkappa_{np}^{(\lambda)}$ (see ref. [35]) can be easily rewritten using the isoscalar $\varkappa_0^{(\lambda)}$ and isovector $\varkappa_1^{(\lambda)}$ constants ($\lambda \neq 1$)

$$(\varkappa_0^{(\lambda)} + \varkappa_1^{(\lambda)})(X_n^i + X_p^i) - 4\varkappa_0^{(\lambda)}\varkappa_1^{(\lambda)} X_n^i X_p^i = 1 \;, \qquad (15)$$

where

$$X_p^i = 2\sum_{\nu\nu'} \frac{f(\nu\nu')\tilde{f}(\nu\nu')\,\mathcal{U}^2_{\nu\nu'}(\varepsilon(\nu)+\varepsilon(\nu'))}{(\varepsilon(\nu)+\varepsilon(\nu'))^2-\omega_i^2}$$

the notations are the same as in refs. [2,36].

Instead of calculating for each state i the reduced probabilities of $E\lambda$-transitions

$$B(E\lambda;\,0^+0 \rightarrow I_i^{\pi_i}K_i) = (00\lambda\mu \mid I_i K_i)^2 M^2(\omega) \qquad (16)$$

and summing them in a certain energy interval, we use the method of the direct calculation of averaged characteristics. We introduce the strength function

$$b(E\lambda,\omega) = (00\lambda\mu \mid IK)^2 \sum_i M^2(\omega_i)\rho(\omega-\omega_i)\,, \qquad (17)$$

$$\int_{\omega}^{\omega+\Delta\omega} b(E\lambda,\omega')d\omega' \approx \sum_{i\,\Delta\omega} B(E\lambda;\,0^+0 \rightarrow I_i^{\pi_i}K_i)\,, \qquad (17')$$

which has the following form for λ =2

$$B(E2,\omega)=\langle 00\lambda\mu|IK\rangle^2\,\frac{2-\delta_{\lambda 0}}{\pi}\,e^2$$

$$\mathrm{Im}\left\{\left[\frac{(1+e_p^{(2)})^2X_p+(e_n^{(2)})^2X_n-\{\kappa_0^{(2)}(1+e_p^{(2)}-e_n^{(2)})^2+\kappa_1^{(2)}(1+e_p^{(2)}+e_n^{(2)})^2\}X_nX_p}{1-(\kappa_0^{(2)}+\kappa_1^{(2)})(X_n+X_p)+4\kappa_0^{(2)}\kappa_1^{(2)}X_nX_p}\right]_{\omega\to\omega+\frac{i}{2}\Delta}\right\} \quad (18)$$

where $e_n^{(2)}$, $e_p^{(2)}$ are the effective charges /36,37/.

Now we consider the giant quadrupole resonances. The strength functions $B(E2,\omega)$ for E 2-transitions to the states with $I^\pi=2^+$ and K =0,1,2 in ^{150}Nd are given in fig.4. The constant $\kappa_0^{(2)}$ was fixed earlier in the description of the low-lying states, $\kappa_1^{(2)}=-3\kappa_0^{(2)}$, $e_n^{(2)}=e_p^{(2)}=0$, Δ =0.4 MeV.

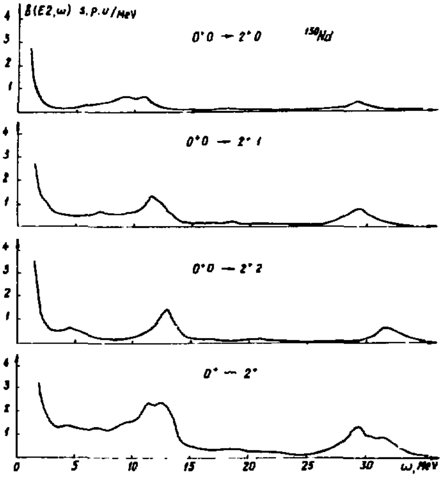

Fig.4. Strength functions B (E 2, ω)(s.p.u./MeV) in ^{150}Nd for E 2 transitions from the ground state to $I^\pi K$ =2^+0,2^+1, 2^+2 and the total value.

in deformed nuclei the giant quadrupole resonances comprise the mixture of components with K =0,1 and 2. The positions of maxima of the resonances with K =0,1,2 do not coincide, and this results in the broadening of isoscaler and isovector resonances in deformed nuclei as compared to spherical nuclei. This is consistent with experimental data[38-40]. Note, that the width we have calculate of the isoscalar maxima with K =0,1 and 2 is essentially different from the calculation results in ref.[39] for ^{154}Sm and qualitatively agree with the calculation results in ref.[41] for ^{160}Dy

Table 3 shows the energies and widths of giant quadrupole and octupole resonances calculated in refs.[36,37] for a number of doubly even deformed nuclei. It is seen from Table 3 that in nuclei of the rare-earth region the energy of the isoscalar resonance is about 12 MeV and of the isovector one about 30 MeV. In the region of actinides the isoscalar resonance is located at energy 10 MeV, the isovector at energy 28 MeV. The calculation results of the energies of isoscalar resonances agree with the available experimental data[38-40,42]. The position of the isovector resonance is determined by the quantity $\kappa_1^{(1)}$. If one assumes $\kappa_1^{(1)} = -1.5\,\kappa_1^{(0)}$ then the energy of the isovector resonance for nuclei of the rare-earth region will be 25 MeV and in the region of actinides 22 MeV what is close to the assumed measured value[43]. Note, that when solving equation (15) we take into account the isoscalar and isovector excitation modes simultaneously, thus, the notions isoscalar and isovector resonances are conventional. It is seen from the Table that the position of giant quadrupole resonances slightly differs inside the rare earth region and that of actinides. This change is even

Table 3.

Energies and widths of isoscalar and isovector quadrupole and octupole giant resonances in deformed nuclei

Nuclei	λ = 2				λ = 3			
	T = 0		T = 1		T = 0		T = 1	
	$\bar{E}$, MeV	$\bar{\Gamma}$, MeV	$\bar{E}$, MeV	$\bar{\Gamma}$, MeV	$\bar{E}$, MeV	$\bar{\Gamma}$, MeV	$\bar{E}$, MeV	$\bar{\Gamma}$, MeV
[illegible]Nd	12.2	2.7	30.3	3.6	20.5	7.5	42.0	2.5
^{152}Sm	12.3	2.4	30.5	4.0	-	-	-	-
^{154}Sm	12.2	2.5	30.5	3.5	21.0	5.0	42.0	3.5
^{154}Gd	12.3	2.5	30.0	3.5	21.0	6.0	42.0	2.5
^{156}Gd	12.3	2.5	30.5	3.6	-	-	-	-
^{158}Gd	12.4	2.5	30.5	4.2	-	-	-	-
^{160}Dy	12.5	2.0	30.2	4.0	-	-	-	-
^{162}Dy	11.[illegible]	1.9	30.5	4.6	20.0	5.5	40.5	3.0
^{164}Dy	11.8	1.5	30.5	4.9	-	-	-	-
^{166}Er	11.9	1.5	30.5	5.3	20.0	5.0	40.5	2.0
^{168}Er	12.0	1.5	30.7	5.5	-	-	-	-
^{172}Yb	11.8	1.4	30.5	5.9	19.5	4.5	40.5	2.5
^{174}Yb	11.8	1.3	29.5	5.4	-	-	-	-
^{176}Yb	11.9	1.4	30.5	5.8	20.0	4.5	40.5	2.0
^{178}Hf	11.7	1.5	30.0	5.8	-	-	-	-
^{230}Th	9.8	1.6	28.6	2.7	17.5	3.5	37.0	3.5
^{232}Th	9.8	1.8	28.4	2.7	17.5	3.5	37.0	3.5
^{234}U	9.5	2.0	28.2	2.9	17.0	2.5	37.5	5.0
^{236}U	9.5	1.8	28.5	3.0	-	-	-	-
^{238}U	9.6	1.5	28.4	3.2	17.5	3.0	37.0	2.5
^{242}Pu	9.8	1.5	28.5	3.6	-	-	-	-
^{244}Cm	9.5	1.5	27.5	4.3	17.0	3.0	36.0	2.5

less than the change of the first $K^{\pi} = 2^{+}$ states. The change of energies and B ($E\lambda$)-quantities for the first one-phonon states is connected with the charge of the chemical potential in the transition of one nucleus to another and, thus, with the charge of the matrix elements corresponding to the first poles. The position of giant resonances is determined by the matrix elements with $\Delta N = \lambda$ and by the corresponding poles the role of which does not change in the transition of one nucleus to another. This fact accounts for a slight charge of the energies of giant quadrupole and octupole resonances in the transition from one nucleus to another.

Now we consider the giant octupole resonances. The strength functions b ($E3$, ω) for $E3$ -transitions to the states $I^{\pi}=3^{-}$ and K =0,1,2,3 in ^{238}U are given in fig.5. The bottom figure represents the total value for $\varkappa_1^{(3)}$ = -4.5 $\varkappa_0^{(3)}$ and $\varkappa_1^{(3)}$=0 calculated with $\varepsilon_n^{(3)} = \varepsilon_p^{(3)} = 0$, Δ =0.4 MeV. It is seen from the figure that the introduction of $\varkappa_1^{(3)} \neq 0$ results both in the formation of the isovector resonance and in the shrinpage of the isoscalar resonance. If we assume $\varkappa_1^{(3)}$=-1.5 $\varkappa_0^{(3)}$, then the energy of the isovector octupole resonance will decrease by 9 MeV.

Let us turn to the model independent energy weighted sum rule (EWSR) and write it down for $\lambda \neq 1$ in the form:

$$\sum_i \omega_i B(E\lambda, \omega_i) = 4{,}8\, \lambda (3+\lambda)^2 \frac{Z}{A^{2/3}} B(E\lambda)_{sp}\ MeV \qquad (19)$$

Table 4 represents the quantities of EWSR for each value of K for the quadrupole and octupole resonances. It is seen that the contribution of all K is considerable and they should be taken into account simultaneously.

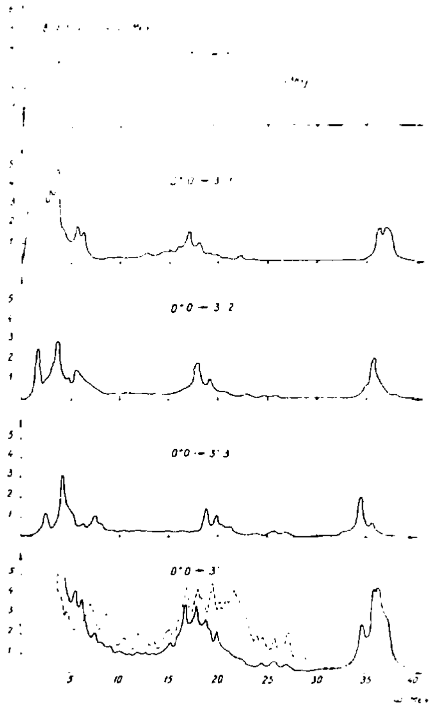

Fig. 5. Strength function $b(E3,\omega)$(s.p.u./MeV) in ^{238}U for $E3$-transitions from the ground state to the states $I^{\pi}K = 3^-0, 3^-1, 3^-2, 3^-3$ and the total value. In the bottom figure the dashed like represents the total value $b(E3,\omega)$ calculated with $\varkappa_1^{(3)} = 0$.

From eq.(19) it follows that the model independent energy weighted sum rule has the following values; for ^{150}Nd $\lambda = 2$ 492, $\lambda = 3$ - 1062; for ^{238}U $\lambda = 2$ - 555, $\lambda = 3$ - 1197. The table shows that the model independent EWSR is exhausted by 90%.

This gives evidence that in the calculations [36,37] almost all the required part of the configurational space has taken into account.

Table 4

Calculated values EWSR in ^{150}Nd and ^{236}U

Nucleus	$\sum_i \omega_i B(E\lambda;\omega_i)$ (S P U MeV)										
	$\lambda=2$					$\lambda=3$					
	$K=0$	$K=1$	$K=2$	all K	K%	$K=0$	$K=1$	$K=2$	$K=3$	all K	K%
^{150}Nd	90.5	199	158	448	88	141	283	242	287	853	78
^{236}U	112	246	202	560	98	177	355	301	251	1084	87

Table 5

Calculated values of EWSR for $\lambda=2$ in ^{150}Nd

	$\sum_i \omega_i B(E3;\omega_i)$ (s.p.u. MeV) Energy intervals, MeV				
	0–2	2–10	10–14	14–26	26–34
$K=0$	5.1	21.9	13.6	15.0	33.1
$K=1$	21.5	30.8	39.0	31.0	73.0
$K=2$	10.6	16.0	41.2	21.6	61.0
all K	37.2	69	94	67	168

Tables 5 and 6 represent the EWSR -value in ^{150}Nd for $\lambda=2$ and in ^{236}U for $\lambda=3$ in different energy intervals. It is seen from these tables that in the region of the isoscalar quadrupole resonance 21% of EWSR is concentrated, and in that of the isoscalar octupole resonance 17% of EWSR. In the region of the isovector quadrupole resonance 38% of EWSR is concentrated, and that of the isovector octupole resonance about 35% of EWSR. It is difficult to determine the resonance widths due to a complex dependence of $\mathcal{E}(E\lambda, \omega)$ on ω what is seen from figs.5 and 6.

Table 6

$\sum \omega_i B(E3\, \omega_i)$ (s.p.u. MeV) Energy intervals, MeV					
	0-2	2-16	16-20	20-34	34-38
K =0	16	39	27	18	65
K =1	25	89	51	36	128
K =2	3,5	65	62	45	108
K =3	0.9	54	43	66	77
all K	46	247	183	165	378

Table 7

Characteristics of the isoscalar quadrupole resonance in ^{165}Ho

	Exp. /45/	Calc./47/
Energy, MeV	11.6 ± 0.2	12.5
width, MeV	3.6	3.5
contr.in EWSR,%	21 ± 4	25
B(E2) e^2 barn^2	0.15 ± 0.03	0.14

Table 3 shows the average energy intervals of the location of giant multipole resonances. Note, that the calculated widths of the isoscalar quadrupole resonances are somewhat smaller than the experimental ones./38-40/

In ref./44/ the mathematical apparatus for calculating $\mathcal{E}(E\lambda)$ -quantities in odd-A deformed nuclei is developed and the giant quadrupole isoscalar resonance is calculated.

Table 7 represents part of these results. It is seen that a good description of the resonance position and width is obtained for [illegible]He .

Our calculations of giant dipole resonances give rather a good description of the available experimental data [31,46].

In the calculation of giant resonances, one should consider the fragmentation of one-phonon states. The widths of giant resonances should be affected by the admixture of two-phonon components comprising the low-lying and high-lying collective phonons to the one-phonon component. That is the calculations should be performed with the wave function (2). This is most important for spherical nuclei in which first calculations are being already performed [47].

Thus, within the framework of our model rather a good description of the energies of giant multipole resonances is obtained, and the apparatus for describing their widths and fine structure is developed.

7. In the framework of the general semimicroscopic approach developed in refs. [5,48,49], the wave function of the highly excited state of an atomic nucleus is represented as the expansion in number of quasiparticles. The neutron, radiational and α - widths are expressed through the coefficients of this wave function. As a rule, only few-quasiparticle components of the wave-functions are engaged in the experiments in which the excitation and decay of individual intermediate and high excitation states are studied. So, when in the study of the neutron resonances, there appear few-quasi-particle components comprising 10^{-3}-10^{-8} part the normalization of the wave functions of neutron resonances. One should note, that only for these few-quasiparticle com-

ponents of the wave functions, the regularities of the nuclear statistical model are valid.

In the framework of this approach based on the operator form of the wave function of a highly excited state, one can make general conclusions on the properties of these states, for instance, on the value of the magnetic moments[50], on the correlation between neutrons and radiational widths[51] on the peculiarities of α -decays[48,52] and so on. It is shown that in deformed nuclei due to richness of the collective excitation modes at the energy $\mathcal{E} = B_n$, the wave functions are so complex that the nonstatistical effects connected with the few-quasiparticle components appear in very rare cases. Thus, one may see the correlations for E 1-transitions with I^- of s -wave resonances to the ground and first 2^+ rotational states in ^{158}Gd , ^{160}Dy , ^{172}Yb , ^{184}W as well as the correlations between the neutron and E 1-radiational widths for the transitions to the states with $K^\pi = 2^+$ from the 3^- resonance in the reaction $^{173}Yb(n,\gamma)^{174}Yb$. There are experimental indications to the existence of such correlations[53].

Based on the general approach, one may find in which nuclei the neutron valence model is valid or not. The analysis performed in refs.[15,54] for the molybdenum isotopes has shown that the neutron valence model should work well in ^{99}Mo and considerably worse in ^{93}Mo . This is due to the fact that near B_n there are no three-quasiparticle states from which the E 1-transitions may proceed to the low-lying states in ^{99}Mo , but there are such states in ^{93}Mo. It is known[55] that there are deviations from the neutron valence model in ^{93}Mo . Recent experiments[56] have confirmed our prediction about the validity of the neutron [illegible] for [illegible] M_1. The general approach allows one to analyse [illegible] peculiarities of the radiational strength functions. Such [illegible] was performed in ref.[57] for [illegible]. At present there [illegible] for the calculation of radiational strength functions [illegible] nuclei [illegible] within the framework of our model.

The study of the fragmentation of two-, three-, and four-quasiparticle states is of most interest. It is important to study experimentally the two-, three-, and four- nucleon transfer reactions. Much attention, for instance, should be paid to the measurement of neutron and partial radiational widths for the neutron resonances in the reactions $^{176}Lu(n,\gamma)^{177}Lu$ and $^{177}Lu(n,\gamma)^{178}Lu$, if one uses as a target the three-quasiparticle isomer 23/2$^-$ in ^{177}Lu (see ref.[5]). The importance of the use of unstable targets in the neutron spectroscopy was mentioned in ref.[4]; we should support such measurements in every possible way (the first of them have already appeared [illegible]).

Up to now only the processes connected with few-quasiparticle components of the wave functions of intermediate and high excitation energies were discussed. In ref.[60] the problem of the value of many-quasiparticle components of the wave functions is [illegible]stated. The assumption was also expressed that the wave functions of the intermediate excitation energy and neutron resonances have rather a large many-quasiparticle components. This is due to the fact that the interactions between quasiparticles and the quasiparticle-phonon interaction at these energies can not fragmentate the many-particle states so strongly as the single-particle ones.

The ways of the experimental detection of large many-quasiparticle components of the wave functions of neutron resonances were discussed in refs.[15,49,60]. Now the most available way of clarifying the role of many-quasiparticle components is the study of E 1-, M 1- and E 2 transitions from the states of intermediate and high excitation energies to the states with the energy by (1.0-1.5) MeV less than their. Possibly, the probabilities of such γ -transitions may be evaluated in the study of the subsequent α -decay of an excited state, the fission or neutron emission. The observation of γ -transition cascades, which reduced probabilities **related to the single-particle ones**, gives evidence to the existence of large many-quasiparticle components in the wave functions of neutron resonances and in the states of the intermediate excitation energy. The information about the values of individual four- and six- quasiparticle components can be **obtained** in the study of the γ transition from neutron resonances to the states of the intermediate excitation energy. For instance, in ref.[61] the relative intensities of a number of γ -transitions between the high-spin levels with energies up to 5.6 MeV in ^{95}Tc were measured. Possible ways of detection of large many-quasi-particle components of the wave functions of intermediate and high excitation energies are schematically presented in fig.6.

The most promising method of measuring of the values of the largest components of neutron resonance wave functions is the study of the reaction (n , γ , α) with the subsequent evaluation of intensities of γ -transitions between the neutron resonances and states lying by (1-2) MeV lower. The experimental results of Popov and collaborators[62] show that there are relatively large components in the wave functions of these states.

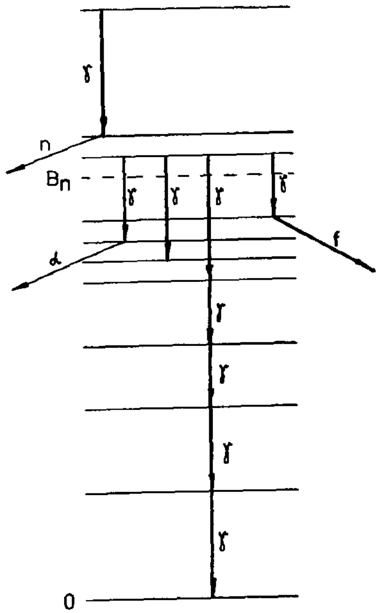

Fig.6. Symbolic representation of the ways of experimental detection of large many quasiparticle components in the wave functions of highly excited states.

It is stated in ref./63/ that the semimicroscopic nuclear theory may serve as a base for a unified descritpion of low, intermediate and high excitation states of complex nuclei. In the present talk we have presented the results of such a unified descritpion. It has been shown that:

i) A satisfactory description of each level has been obtained for the low-lying states;

ii) The description of the fragmentation of few-quasiparticle states is given for the intermediate energy states. For the experimental study of the fragmentation it is neccessary to measure the one-nucleon transfer reactions with fixed momentum transferred;

iii) The description of S -, p - and d -wave neutron strength functions is given for the neutron resonances, and the method of calculating the radiational strength functions is developed;

iv) The description of the energies of giant multipole resonances is obtained, and the method of calculating their widths and fine structure is developed;

v) Due to the fact that the main part of the experimental information on the states of intermediate and high excitation energy concerns few-quasiparticle components of their wave functions it is necessary to find the values of many-quasiparticle components.

One may conclude that there is a model for the description of few-quasiparticle components of the states at low, intermediate and high excitation energies. The few-quasiparticle components of the states at intermediate and high excitation energies are represented to be averaged in some energy interval, i.e., as the correspondding strength functions.

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