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A NEW APPROACH TO THE STUDY OF THE STRUCTURE OF NEUTRON RESONANCES

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I. Introduction

A series of papers devoted to a new approach to the study of the structure of highly excited states of atomic nuclei has for an object firstly to clarify how proceeds the complication of the state structure with increasing excitation energy, and secondly to develop a unique description for low-lying states, states of intermediate excitation energy and highly excited states. The realization of this aim proceeds along two lines; a general semi-microscopic description on the basis of the operator form of the wave functions of highly excited states and numerical calculations on the basis of the model taking into account the interactions of quasiparticles with phonons.

2. General Semi-Microscopic Approach

We give the basic formulas of the general semi-microscopic description of the structure of highly excited states developed in refs. /1-4/.

The wave function of a highly excited state is represented as a sum of terms, each containing the operators of quasiparticles and nucleons for the neutron and the proton systems. With increasing excitation energy a complication of the structure of nuclear states is seen from the fact that the wave function contains components with ever-increasing number of quasiparticles.

Let us construct, for example, the wave function of the highly excited state of an odd-N spherical nucleus. It consists of an one-quasiparticle neutron component (n), three-quasiparticle components ($n2p$), (3n) (one neutron and two proton quasiparticles and three neutron quasiparticles) and five-quasiparticle components (n4p), (3n, 2p), (5n) etc. We write it in the following form:

$$
\begin{aligned}
\Psi_\lambda(I^\pi M) = {} & b_I^{\lambda n}(j^n)\,\delta_{I,j^n}\,\delta_{M,m^n}\,\alpha^+_{j^n m^n}\,\Psi_0 + \\
& + \sum_{\substack{j_1^n, j_2 j_3 \\ m_1^n, m_2, m_3}} \sum_t b_I^{\lambda n 2t}(j_1^n m_1^n; j_2 m_2 j_3 m_3)\,\alpha^+_{j^n m^n}\alpha^+_{j_2 m_2}\alpha^+_{j_3 m_3}\,\Psi_0 + \\
& + \sum_{\substack{j_1^n, j_2, j_3, j_4, j_5 \\ m_1^n, m_2, m_3, m_4, m_5}} \sum_{t,t'} b_I^{\lambda n 2t 2t'}(j_1^n m_1^n, j_2 m_2, j_3 m_3, j_4 m_4, j_5 m_5)\,\alpha^+_{j_1^n m_1}\alpha^+_{j_2 m_2}\alpha^+_{j_3 m_3}\alpha^+_{j_4 m_4}\alpha^+_{j_5 m_5}\,\Psi_0 + \\
& + \cdots + \sum_{t\,\eta} b_I^{\lambda n \Omega_\eta(t)}(j^n)\,\delta_{I j^n}\,\delta_{M, m^n}\,\alpha^+_{j^n m^n}\,\Omega^+_\eta(t)\,\Psi_0 + \cdots,
\end{aligned}
\tag{1}
$$

where α^+_{jm} is the quasiparticle creation operator (see ref. [5]). The coefficients b^λ define the contribution of the corresponding quasiparticle component, λ is the number of an excited state with given $I^\pi M$. The index $t=n$ indicates the neutron and $t=p$ the proton systems. In the wave function (I) the products $(\alpha^+_{jm}\alpha^+_{j-m})_{I=0}$ are replaced by the phonon pairing vibration operators $\Omega^+_\eta(t)$ the explicit form of which is given in refs. [4,5]. By Ψ_0 we mean the product of quasiparticle or phonon vacua for the neutron and proton systems.

Besides the operators $\Omega^+_\xi(t)$, the wave function (I) contains no other phonon operators. As far as the phonon operators are written in the form of a superposition of two-quasiparticle operators, the corresponding terms are thought of as being already included in the wave function (I). If needed, it is possible to include explicitly the operators of quadrupole, octupole and other phonons in the wave function (I).

It is unknown how proceeds the rotation of a nucleus in the highly excited state. Thus, it may be expected that the projection K of the total angular momentum on the nuclear symmetry axis is not a sufficiently good quantum number and the wave functions of the highly excited states of deformed nuclei are a superposition of terms with different K . The wave function of an odd-N deformed nucleus is written in the form

$$\Psi_\lambda(I^\pi M) = \sqrt{\frac{2I+1}{16\pi^2}} \sum_{K>0} d_I^\lambda(K)\left\{D^I_{MK}(\theta_e)\Psi_\lambda(K^\pi) + (-1)^{I+K} D_{M-K}(\theta_e)\Psi_\lambda(-K^\pi)\right\}, \tag{2}$$

$$\Psi_\lambda(K^\pi) = \sum_s b_K^{\lambda n}(s)\alpha^+_{s+}\Psi_0 + \sum_{\substack{s_1,q_2,q_3\\ \sigma_1,\sigma_2,\sigma_3}} \sum_t b_K^{\lambda n 2t}(s_1\sigma_1, q_2\sigma_2, q_3\sigma_3)\alpha^+_{s_1\sigma_1}\alpha^+_{q_2\sigma_2}\alpha^+_{q_3\sigma_3}\Psi_0 + \tag{3}$$

$$+ \cdots + \sum_s \sum_{t\xi} b^{\lambda n \Omega_\xi(t)}(s)\alpha^+_{s+}\Omega^+_\xi(t)\Psi_0 + \cdots \quad .$$

The set of quantum numbers defining the single-particle state is denoted by ($s\sigma$) for the neutron system and by ($q\sigma$) for both the neutron and proton ones. In each term of eq.(3), the summation over σ is subjected to an additional condidition of the type $K=\sigma_1 K_1+\sigma_2 K_2+K_3\sigma_3$. The wave functions of the highly excited states of even spherical and deformed nuclei are given in ref.[4]

It should be noted that the operator form of the highly excited state wave function written in the form of (I) or(3) is not the most general one. When constructing the wave function an approximation is used which consists in that in each term the operators of quasiparticles and pairing vibration phonons act on the wave funotions which are either a quasiparticle or a phonon vacuum.

It is worth noting that if the wave function describes a state the energy of which exceeds the neutron binding energy then, strictly speaking, it should be written as

$$\Psi_\lambda(\mathcal{E})=\sqrt{\frac{1}{2\pi}\,\frac{\Gamma_{n\lambda}}{(\mathcal{E}-\mathcal{E}_\lambda)^2+\Gamma_n^2/4}}\;\Psi_\lambda(I^\pi M)\,, \qquad (4)$$

where $\mathcal{E}_\lambda$ is the resonance energy. This corresponds to the fact that a neutron can be emitted, i.e. when a neutron interacts with the target-nucleus A-1 the elastic scattering channel is open. However, this energy factor is, as a rule, unimportant when calculating transitions from highly excited states.

At excitation energies close to the neutron binding energy B_n or higher the wave functions contain thousands of various few- and many-quasiparticle components. Such wave functions possess the properties of the compound states introduced by N.Bohr. In fact, the formation of a highly excited state can proceed through one components, and the decay through other components of the wave function. Therefore, in many cases, the main requirement that the decay of compound states is independent of the mode of their formation will be fulfilled. Since the wave function contains a variety of components some few-quasiparticle components must have, as a rule, small values. This will lead to a significant hindrance of the probabilities of gamma transitions to low-lying states. Therefore, the half-life of a highly excited state must be much longer than that of a one- or two-quasiparticle state.

These wave functions may be used for describing excited states from (2-3) MeV to such energies at which resonances do not overlap yet, i.e. when the condition

$$\Gamma_n \ll D \tag{5}$$

holds. This implies that the neutron width Γ_n is much smaller than the mean spacing D between levels with given I^{π}.

Using the operator form of the wave functions (I) and (3), in refs. /1-4/ the reduced neutron $\Gamma^{\circ}_{n\lambda}$, radiative $\Gamma^{\circ}_{\gamma\lambda f}$ and alpha $\Gamma^{\circ}_{\alpha\lambda f'}$ widths, as well as the strength functions for s- and p- wave neutrons are expressed in terms of the coefficients b^{λ}. If there is the operator form of the wave function of a neutron resonance then in the language of the wave function coefficients it is possible to describe the neutron entrance channels and obtain the same relations as in models of valency-neutron and doorway states. In calculations performed in the framework of the doorway state model, only a part of the components of the wave function (I) or(3) is used, and the connection with the remaining components is mentioned. The main difference of our approach from the doorway state model consists in that we take the wave function as a superposition of components with different number of quasiparticles, while in the doorway state model the wave function contains a definite number of quasiparticles.

We consider the reduced neutron widths $\Gamma^{o}_{n\lambda}$. The reduced width amplitude for the state $|\lambda\rangle$ in the channel c is usually[6] written as

$$\gamma_{\lambda c} = \sum_{S} \langle \lambda | \varphi_S \Psi_c \rangle \gamma_S ,$$

where $\Psi_c \varphi_S$ is the product of the internal wave functions of the nucleus in the channel C and of the one-particle state, γ_S is the one-particle amplitude for the S -state. In our treatment, the quantity $\langle \lambda | \varphi_S \Psi_c \rangle$ can be replaced by the matrix element $(\Psi^{*}_{\lambda}(I^{\pi}) a^{+}_{jm} \Psi_i)$, where Ψ_i is the wave function of the ground state of a target-nucleus. If we take as the target an even-even spherical nucleus and Ψ_{λ} in the form (I) then we get

$$(\Psi^{*}_{\lambda}(I^{\pi}) a^{+}_{jm} \Psi_c) = b^{*\lambda n}_{I}(j) u_j .$$

The reduced neutron width can be written as

$$\Gamma^{o}_{n\lambda} = \Gamma^{o}_{s.p.} | b^{\lambda n}_{I}(j) u_j |^2 + \Gamma^{oo}_{n\lambda} , \qquad (6)$$

where $\Gamma^{o}_{s.p.}$ is the single-particle value of the neutron width, the function u_j indicates that the states j must be particle ones. The term $\Gamma^{oo}_{n\lambda}$ is responsible for a more complicated neutron capture mechanism. It may contain the same terms as in the doorway state model (see refs.[7,8]). If as a doorway state we take three-quasiparticle components of the wave function (I), then $\Gamma_{n\lambda}^{cc}$ will contain components $b_I^{\lambda n 2t}(j_1 m_1, j_2 m_2, j_3 m_3)$ and $b_I^{\lambda n \Omega_3 (t)}(j)$.

We give an expression for the reduced radiative width which is connected with the $E\lambda$ transition matrix element as follows:

$$\Gamma^{\circ}_{\gamma\lambda f} = \frac{1}{2I+1}\,\frac{1}{2\lambda+1}\Big|\sum_{M,m_f}(j_f m_f \lambda_M | IM) M(E\lambda; I^{\pi}\lambda \to j_f m_f)\Big|^2 . \qquad (7)$$

The matrix element for the $E\lambda$ transition from the highly excited state described by the wave function (I) to the one-quasiparticle state $\alpha^{+}_{j_f m_f}\Psi_0$ is

$$\begin{aligned} M(E\lambda; I^{\pi}\lambda \to j_f m_f) = {} & \frac{(-1)^{\lambda}}{\sqrt{2\lambda+1}}(-1)^{j_f - M}(j_f m_f I -M|\lambda\mu) b_I^{\lambda n}(j)\delta_{I,j} v^{(-)}_{jj_f}\langle j|\Gamma(E\lambda)|j_f\rangle - \\ & -\frac{1}{\sqrt{2\lambda+1}}\sum_{\substack{j_1,j_2\\ m_1,m_2}}(-1)^{j_1+j_2-\mu}(j_1 m_1 j_2 m_2|\lambda\mu)\left[b_{IM}^{\lambda n 2p}(j_f m_f, j_1 m_1, j_2 m_2) + 3b_{IM}^{\lambda n 2n}(j_f m_f j_1 m_1, j_2 m_2)\right]\cdot \\ & \cdot u^{(+)}_{j_1 j_2}\langle j_2|\Gamma(E\lambda)|j_1\rangle - \frac{1+(-1)^{\lambda}}{4\sqrt{2\lambda+1}}(-1)^{j_f - I - \mu}(j_f - m_f IM|\lambda -\mu) b_I^{\lambda n \Omega_3 (n)}(j)\delta_{j,I} u^{(+)}_{j_f j}\cdot \\ & \cdot \psi^{\varphi}_{j_f j_f}\langle j|\Gamma(E\lambda)|j_f\rangle + \cdots \end{aligned} \qquad (8)$$

The first term in (8) corresponds to the valency neutron model. The matrix elements for the $E\lambda$ and $M\lambda$ transitions to one-phonon states or to states quasiparticle plus phonon are more complicated.

The matrix elements of alpha transitions from highly excited to low-lying states have a variety of components of the wave functions of highly excited states[1,4].

3. Structure of Neutron Resonances

We use the operator form of the wave functions of neutron resonances in order to clarify their structure. Let us look what information about the b^{λ} coefficients in (I) and (3) can be extracted from neutron resonance experiments.

First of all, we clarify what is the order of magnitude of one-quasiparticle and two-quasiparticle components of the wave functions of neutron resonances. We consider as an example the case when an s-wave neutron is captured by an even-even spherical nucleus. If in eq.(6) we reject the term $\Gamma_{n\lambda}^{oc}$ the reduced neutron width is then defined as

$$\Gamma_{n\lambda}^{o} = \Gamma_{s.p.}^{o} \, | b_I^{\lambda n}(j) \mathcal{U}_j |^2 , \qquad (6')$$

where $j^{n} = 1/2^{+}$. We assume that $\mathcal{U}_j^2 \approx 1$, then

$$\Gamma_{n\lambda}^{o} \approx \Gamma_{s.p.}^{o} \, | \bar{b}^{\lambda} |^2 , \qquad (9)$$

we make an averaging over a number of resonances and get

$$< \Gamma_n^{o} > \approx \Gamma_{s.p.}^{o} \, | \bar{b} |^2 . \qquad (10)$$

Table I gives experimental data on $<\Gamma_n^{o}>$ and the $|\bar{b}|^2$-values. It is seen that $|\bar{b}|^2$ is, in the order of magnitude, 10^{-6} in nuclei around closed shells in which fragmentation is not strong enough, and in far lighter nuclei with A=50–60 where the strength function for S-wave neutrons has a maximum. In nuclei with A $\approx$ 100, $|\bar{b}|^2$ is 10^{-7} – 10^{-8} for one-quasiparticle components $S_{1/2}$ and for those two-quasiparticle components one of which is $S_{1/2}$. In the same nuclei, $|\bar{b}|^2$ is 10^{-6} for one-quasiparticle components $3p_{1/2}$ and $3p_{3/2}$ which is related to the maximum of the strength function for p -wave neutrons.

In strongly deformed nuclei the s -wave strength functions are not small, however, the $|\bar{b}|^2$- values are very small, 10^{-8} - 10^{-9}. This decrease of $|\bar{b}|^2$ is connected with a variety of collective excitation branches in deformed nuclei as compared with spherical nuclei. A large number of collective excitations in deformed nuclei has led to stronger fragmentation of one-particle components over many nuclear levels. The decrease of $\langle \Gamma_n^o \rangle$ in deformed nuclei compared to spherical ones may be partially due to the difference of equilibrium deformations of the highly excited state and the ground state of the target-nucleus.

The magnitude of the one- and two-quasiparticle components of the wave functions of highly excited states estimated from the experimental data on $E1$ transitions to the ground and low-lying states coincides with the estimate obtained from the neutron widths. According to refs./3,4/ for nuclei with A from 50 to 250 the one- and two-quasiparticle components of the wave functions of neutron resonances take the values in the following interval:

$$|\bar{b}|^2 = 10^{-6} \cdot 10^{-9} .$$

In refs./3,4/ there are expressions for the reduced neutron widths in the case when a slow neutron is captured by an odd-odd nucleus. Then it is possible to obtain information on the magnitude of three-quasiparticle components ($\rho 2n$). It follows from the experimental data that the $|\bar{b}|^2$ -values for three-quasiparticle components of states with high spins are smaller than the $|\bar{b}|^2$ -values for one- and two-quasiparticle components. This appears to be due to the fact that the fragmentation of three-quasiparticle components is at the initial stage and the main part of the strength of three-particle states is concentrated on a few levels. Evidence on the magnitude of three- and four-quasiparticle components of the wave functions of neutron resonances will be richer if in neutron-spectroscopic studies unstable nuclei and isomers are used as targets/3/.

The role of many-quasiparticle components of the wave functions of neutron resonances is not clear. The fact that some few-quasiparticle components give a contribution of the order 10^{-6} - 10^{-9} to the wave function normalization does not mean at all that the number of the wave function components is many millions and that all of them are small. It is quite possible that one or several many-quasiparticle components give a predominant contribution to the normalization of the highly excited state wave function.

We study what indirect information about the magnitude of four-, five-, six-, seven- and eight-quasiparticle components of the wave functions can be extracted from an analysis of alpha and gamma decay of resonances. It is obvious that the role of many-quasiparticle components must be displayed most strongly in those resonances for which the corresponding strength function is minimum.

We consider $E1$ transitions from highly excited states. Let a slow neutron be captured by an odd-N spherical nucleus and resonances with $I^{\pi} = 1^-$ be produced. The $E1$ transition to the two-phonon 0^+ state involves, in addition to some two-quasiparticle components, which take part in transition to the ground 0^+_1 state, a large number of four- and six-quasiparticle components of the highly excited state wave functions. $E1$ transitions to one-phonon 2^+ states involve two- and four-quasiparticle components while $E1$ transitions to two-phonon 2^+ states involve also six-quasiparticle components. If four-quasiparticle components are not small one may expect an increase of the reduced probabilities of $E1$ transitions to one-phonon states compared with $E1$ transitions to the ground states. An increase of the reduced probabilities for $E1$ transitions to two-phonon states compared with $E1$ transitions to the ground and one-phonon states will give evidence for a noticeable contribution of six-quasiparticle components to the normalization of the highly excited state wave function.

Some information about the total contribution of six- and eight-quasiparticle components can be extracted from experimental data on reduced partial probabilities of alpha decay of neutron resonances. The alpha decay of resonances to the ground states of even-even nuclei involve two- and four- quasiparticle components of the type ($2n$ $2p$). The alpha decay of resonances to one-phonon states involve also a large number of four- and six-quasiparticle components. Thus, if four- and six-quasiparticle components are important, the reduced probabilities for alpha transitions to one-phonon states must be larger than to the ground states. It follows from the experimental data/9/ that when averaging over eight $I^{\pi}=3^{-}$ resonances in ^{148}Sm the reduced probabilities for alpha transitions to the one-phonon 2^{+} state in ^{144}Nd are twice as large as those for alpha transitions to the ground state. These data point to the important role of four- and six-quasiparticle components in the wave functions of highly excited states.

The indirect data on the total contribution of many-quasiparticle components do not answer the question as to whether or not the wave functions of neutron resonances contain large many-quasiparticle components. It seems that this question may be answered by studying $E1$ and $E2$ transitions from neutron resonances to the states the energy of which is lower by (I.0-1.5) MeV than the neutron resonance energy. If, for example, the wave function of a given resonance has a large seven-quasiparticle component then $E1$ and $E2$ transitions should occur from it to states containing appreciable five-quasiparticle components. The reduced probabilities of these transitions must be much larger than for transitions to low-lying states. Observation of large reduced probabilities for gamma transitions from neutron resonances to some states with somewhat lower excitation energy will testify to the presence of large many-quasiparticle components in the wave functions of neutron resonances.

Experimental data on many-quasiparticle components may apparently be obtained from many-particle transfer reaction studies.

As a first step toward clarifying the rotational properties of highly excited states in ref.[10] it is suggested to study how good for them is the quantum number K . The gamma transitions from resonances are assumed to be K- forbidden if for them the condition

$$|I_\lambda - K_f - \lambda| = \nu > 0, \tag{11}$$

holds, where λ is the multipolarity of a gamma transition. The role of the quantum number K in highly excited states may be judged by the rate of hindrance of K -forbidden gamma transitions from resonances compared with K allowed transitions. It may be expected that the role of the quantum number K will be different in different terms of the wave functions of higly excited states. The available experimental data indicate that the hindrance of K -forbidden gamma transitions may be small.

In refs.[4,11] an expression for the magnetic moment of a highly excited state is derived and it is shown that because it contains all the components of the wave function

the magnitude of the magnetic moments for highly excited states must be equal to that for low-lying states. At present there are experimental data[12,13] on magnetic moments of two resonances in ^{168}Er which confirm qualitatively this conclusion.

In refs.[3,10,14,15] on the basis of the semi-microscopic approach, correlations between neutron radiative and alpha widths in neutron resonances are considered. It is shown that correlations can occur when quasiparticle selection rules are valid. The cases are indicated which are the most favourable for correlations to occur.

4. A Model for Describing the Structure of Highly Excited States

The general semi-microscopic approach to the study of the structure of highly excited states should be completed by model calculations. In refs. [2,16] the study is performed by the model in which quasiparticle-phonon interactions are taken into account. We give here the basic assertions and some results obtained in the framework of a modification which is used for studying nonrotational states in odd-A deformed nuclei.

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

$$\begin{aligned} \Psi_i(K^{\pi}) = C_{s_i}^{i} \frac{1}{\sqrt{2}} \sum_{\sigma} \Big\{ \alpha^{+}_{s_i\sigma} &+ \sum_{g,s} D^{gi}_{s_i s\sigma} \alpha^{+}_{s\sigma} Q^{+}_{g} + \\ &+ \sum_{g,g_2} \sum_{s} F^{gg_2 i}_{s_i s\sigma} \alpha^{+}_{s\sigma} Q^{+}_{g} Q^{+}_{g_2} + \\ &+ \frac{1}{\sqrt{3}} \sum_{g,g_2,g_3} \sum_{s} R^{gg_2 g_3 i}_{s_i s\sigma} \alpha^{+}_{s\sigma} Q^{+}_{g} Q^{+}_{g_2} Q^{+}_{g_3} \Big\} \Psi_0 , \end{aligned} \tag{12}$$

where Q_g is the phonon creation operator, $g = \lambda\mu j$, $\lambda\mu$ the multipolarity, j the number of the root of the secular equation for a phonon.

The Hamiltonian of the model contains the average field, interactions leading to superconducting pairing correlations and multipole-multipole forces. The constants of the multipole-multipole interaction are fixed in determining phonons in the appropriate even-even nuclei. The problem is solved as follows. We calculate the average value of the Hamiltonian over the state (12) and on the basis of the variational principle we derive equations for determining the functions $D^{g_i}_{s_0 s\sigma}$ $F^{g g_2 i}_{s_0 s\sigma}$ and $R^{g g_2 g_3 i}_{s_0 s\sigma}$. We reject noncoherent terms in one of the equations and the problem thus reduces to a secular equation of a rather complicated form (see[16]). However, this has resulted in the appearance of extraneous roots.

We have succeeded in overcoming this trouble by taking into account a part of coherent terms.

Numerical calculations are performed for ^{239}U, ^{169}Er, ^{167}Er and ^{165}Dy. In the multipole-multipole interaction phonons with $\lambda = 4$ and $\lambda = 5$ are taken into account in addition to quadrupole with $\lambda = 2$ and octupole $\lambda = 3$ phonons. The replacement of the two-quasiparticle states by the phonons with $\lambda = 4$ and $\lambda = 5$ is explained by the mathematical formalism used. All roots of the secular equations for one-phonon states up to the neutron binding energy B_n are taken into account in the calculations and the one-particle energies and the wave functions of the Saxon-Woods potential are used.

First we perform calculations of the density of states with given K^{π} which is defined by the number of poles

$$\begin{aligned} &\mathcal{E}(s) + \omega_g \\ &\mathcal{E}(s) + \omega_{g_1} + \omega_{g_2} \\ &\mathcal{E}(s) + \omega_{g_1} + \omega_{g_2} + \omega_{g_3} \\ &\dots \end{aligned} \tag{13}$$

This calculation has some advantage compared to the calculation in the framework of the models of independent particles and independent quasiparticles (see ref. /17/) since, in addition to quasiparticles, the phonon states are taken into account.

The calculated average spacing D between the the $1/2^+$ levels as a function of the excitation energy is given in Table 2. It is seen that the density increases strongly with increasing excitation energy. Table 3 gives the average spacing D between the $1/2^+$ levels at excitation energies equal to B_n. It is seen from the table that the calculated density is about half the measured one. Thus a small difference indicates that our model may serve as a good basis for studying highly excited states.

In the case of ^{235}U the number of poles(13)with $K^\pi = 1/2^+$ in the interval 10 keV near B_n is the following: $\mathcal{E}(s) + \omega_g$ - 7 poles; $\mathcal{E}(s) + \omega_g + \omega_{g_2}$:114 poles; $\mathcal{E}(s) + \omega_g + \omega_{g_2} + \omega_{g_3}$: 172 poles; and $\mathcal{E}(s) + \omega_g + \omega_{g_2} + \omega_{g_3} + \omega_{g_4}$:26 poles. That is, the overwhelming majority of poles are five- and seven-quasiparticle ones. Further, when adding phonons with $\lambda = 6$ and $\lambda = 7$ $D = 31.5$ eV, i.e. it leads only to a small increase in the density.

The main merit of the model is the fact that the wave function of a highly excited state has a variety of components. It follows from the calculations for ^{235}U with the wave function (12) which has $R^{g_1 g_2 g_3 i}_{\rho s \sigma} = 0$ that the one-quasiparticle components $(C^i_\rho)^2$ run over the values between 10^{-7} and 10^{-10} which are in agreement with the values of the one-quasiparticle components obtained from neutron widths.

In the previous section the question is as follows. Are there large many-quasiparticle components in the wave functions of neutron resonances? As was already mentioned, at energies close to the neutron binding energy B_n the overwhelming majority of poles (13) are many-quasiparticle ones which points to a possible existence of large many-quasiparticle components in the wave functions. The calculations show that at energies close to B_n the wave functions in ^{239}U contain a number of large components quasiparticle plus two phonons. Interactions between quasiparticles and interactions of quasiparticles with phonons at energies close to the neutron binding energy cannot fragmentize many-particle states as strongly as single-particle states. The calculations performed by means of our model testify in favour of the hypothesis about the presence of large many-quasiparticle components in the wave functions of neutron resonances.

The model considered is more general compared with the models of valency neutron and doorway states, which are particular cases of the former. Indeed, if the neutron width is calculated only with the terms of the wave function (12) containing some components quasiparticle plus phonon then the results for the deformed nucleus are similar to those obtained in ref. /8/ for nuclei around closed shells.

Investigations in the framework of this model are being continued. The composition of the wave functions of highly excited states is being studied, the characteristics of neutron resonances are being calculated etc. A similar model has been formulated for describing highly excited states in odd-A spherical nuclei.

On the basis of the calculations by this model it may be concluded that the mechanism of interactions of quasiparticles with phonons (or interactions of single-particle and collective degrees of freedom) is, to a large extent, resposible for the complication of the structure of states with increasing excitation energy.

5. Complication of the Structure of Nuclear States with Increasing Excitation Energy

Above we have discussed the properties of neutron resonances. Such a particular attention to neutron resonances is connected with a sufficient amount of experimental information for them. It would be much better if it was possible to advance consequtively in the study of the state structure as the excitation energy increases.

We discuss the general picture of the fragmentation of single-particle states over many nuclear levels. We assume that quasiparticle-phonon interactions are the main mechanism which is responsible for fragmentation.

If the energy of a single-particle state is near the Fermi surface energy it may be roughly supposed that 90% of the strength of the single-particle state is concentrated at one nuclear level, (8-9)% - at several low-lying levels and the remaining (1-2)% are distributed over a large number of levels in the interval of 10 MeV and higher. The information about the distribution of the strength of a single-particle state in a wide energy interval is obtained from nonzero strength functions for s -wave neutrons in those nuclei the subshell of which $S_{1/2}$ is the ground or the hole state.

As the single-particle energy moves away from the Fermi surface energy the fragmentation of a single-particle state over many levels becomes stronger. When the single-particle state energy is far away from the Fermi surface by (2-3) MeV about 90% of this strength is distributed over several nuclear levels and the remaining 10% are fragmentized over many levels in a wide energy interval. As the single-particle energy moves away from the Fermi surface energy there proceeds further enhancement of the fragmentation process.

Fragmentation of three-, five- and higher particle states over many nuclear levels occurs. The rate of fragmentation of three-particle states is as strong as the rate of fragmentation of single-particle states at higher excitations etc.

Some approximate conservation laws weaken this rate for one states compared to others. The most striking example is isobaric analogue states which are few-quasiparticle states due to the fact that the conservation law of isotopic spin acts against mixing of states with isotopic spin T_0 to the state with isotopic spin T_0 +1.

Another example are states with large values of spin I in spherical nuclei and quantum number K in deformed nuclei. The wave functions of these states contain no few-quasiparticle components and the fragmentation of many-particle states begins work at higher excitations than that of single-particle states. As a result of both facts, the lowest states of spherical nuclei with large spins I and deformed nuclei with large K are rather pure quasiparticle states.

One should expect that, at still higher excitations, six-, seven- and higher quasiparticle states with large I or K will be observed experimentally. The structure of these states differs strongly from that of other neighboring states.

Since the fragmentation of many-particle states becomes strong at high excitations one may expect rather pure many-quasiparticle highly excited states with small spins.

The rate of fragmentation is different for different nuclei. An especially strong difference is observed in doubly magic nuclei, e.g. in ^{208}Pb and in nuclei which differ from the former by one nucleon. In doubly magic nuclei the first one-phonon states are located more highly than in nonmagic nuclei and the interactions of quasiparticles with phonons are weakened. Therefore the process of fragmentation of one-particle states over many nuclear levels in doubly magic nuclei begins at high energies and develops more slowly than in other nuclei.

In fact, the properties of neutron resonances in ^{207}Pb and ^{209}Pb produced after S-neutron capture differ essentially from the properties of neutron resonances in other nuclei. For example, the average spacing D between the $1/2^+$ levels in ^{207}Pb is larger by about two orders of magnitude than in ^{199}Hg and ^{193}Pt, at the same excitations. Next, in ^{209}Pb a resonance with $\mathcal{E}-B_n$ = 500 keV and anomalously large reduced width $\Gamma^o_{n\lambda}$ = 80 eV is observed. The corresponding one-quasiparticle component $|b^\lambda|^2 = 1.2 \cdot 10^{-4}$ is larger by about $10^3 - 10^4$ than in other nuclei. In ^{207}Pb ten $1/2^+$ states are observed in the interval 200-550 keV the reduced neutron widths of which $\Gamma^o_{n\lambda}$ are large and the corresponding values $|b^\lambda|^2$ reach 2.10^{-5}. The partial radiative widths $\Gamma^o_{\gamma\lambda f}$ for $E1$ transitions from these resonances to the ground state take the values from 1.1 to 10.2 keV.

Thus, the wave functions of nuclei around closed shells have relatively large few-quasiparticle components. Therefore the doorway state model is successfully applied to these nuclei. The behaviour of neutron resonances in nuclei around closed shells must demonstrate the deviation from the regularities of the statistical model which is confirmed experimentally.

On the basis of the information about the components of the neutron resonance wave functions obtained from experimental data and calculations by the model we can represent the structure of states at excitation energies close to the neutron binding energy as follows.

i). One- and two-quasiparticle components have approximately the same value in a large energy interval higher B_n. Their behaviour obeys statistical laws which is a consequence of the developed fragmentation process. The dependence of their values upon A is demonstrated by the strength functions S_0 and S_1 the magnitude of which is defined by the location of the subshells $S_{1/2}$ and $P_{1/2}$, $P_{3/2}$ with respect to B_n.

ii). Three- and four-particle states are fragmentized over many nuclear levels and therefore the three- and four-quasiparticle components of the wave functions must satisfy statistical laws. The exception is light nuclei and nuclei around closed shells (as well as analogue states) where the fragmentation process is not strong enough. This is seen experimentally from the presence of intermediate structures.

iii). For many-quasiparticle components the fragmentation process is at the initial stage and the wave functions may have large many-quasiparticle components and, consequently, individual characteristics. Their characteristics should not apparently obey statistical laws.

These considerations, especially concerning the behaviour of many-quasiparticle components, are in disagreement with the commonly accepted idea about compound states and should be verified experimentally.

We have thus presented the general semi-microscopic approach for studying the structure of highly excited states. On the basis of this approach it is possible to clarify the general regularities and to extract from experimental data information on some components of the wave functions of highly excited states. A detailed theoretical study of the structure of highly excited states can be made using rather simple models.

Thus, the combination of the general semi-microscopic method for treating nuclear states with calculations performed by means of various models can contribute largely to the study of the structure of highly excited states.

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Table I.

Average neutron widths $\langle \bar{\Gamma}_n^0 \rangle$ and one- and two-quasiparticle components $|\bar{b}|^2$.

Nuclei	I^{π}	B_n MeV	$\langle \Gamma_n^0 \rangle$ eV	$\lvert \bar{b} \rvert^2$
^{41}Ca	$1/2^+$	8.364	15	4.10^{-6}
^{51}Cr	$1/2^+$	9.270	8	$2\ 10^{-6}$
^{59}Ni	$1/2^+$	9.003	6	$2\ 10^{-6}$
^{65}Zn	$1/2^+$	7.988	0,4	$1.3\ 10^{-7}$
^{97}Mo	$1/2^+$	6.816	0.1	$1\ 10^{-7}$
^{98}Mo	$2^+,3^+$	8.642	0.004	$0.4\ 10^{-8}$
^{117}Sn	$1/2^+$	6.941	0.012	$1.2\ 10^{-8}$
^{138}Ba	$1^+,2^+$	8.540	0.08	$2\ 10^{-7}$
^{139}Ba	$1/2^+$	4.720	1.8	$2\ 10^{-6}$
^{146}Nd	$3^-,4^-$	7.561	0.004	$0.5\ 10^{-8}$
^{158}Gd	$1^-,2^-$	7.929	0.003	$4\ 10^{-9}$
^{165}Er	$1/2^+$	6.645	0.003	$4\ 10^{-9}$
^{171}Yb	$1/2^+$	6.760	0.003	$4\ 10^{-9}$
^{172}Yb	$0^-,1^-$	8.140	0.001	$1.2\ 10^{-9}$
^{182}Ta	$3^+,4^+$	6.060	0.0005	$0.4\ 10^{-9}$
^{198}Au	$1^+,2^+$	6.497	0.0014	$2\ 10^{-9}$
^{201}Hg	$1/2^+$	6.227	0.2	$0.3\ 10^{-6}$
^{207}Pb	$1/2^+$	6.734	3.4	$5\ 10^{-6}$
^{208}Pb	$0^-,1^-$	7.376	0.6	$1\ 10^{-6}$
^{210}Bi	$4^-,5^-$	4.600	0.13	$0.2\ 10^{-6}$
^{239}U	$1/2^+$	4.800	0.002	$4\ 10^{-9}$

Table 2.

Average spacing D between the $I^{\pi} = 1/2^{+}$ levels in ^{239}U

$\mathcal{E}$, MeV	D, eV	$\mathcal{E}$, MeV	D, eV
1,0	66.7 10^{3}	4.8	33.2
2.0	9.4 10^{3}	5.0	27.0
3.0	1.43 10^{3}	6.0	7.1
4.0	165	8.0	0.5

Table 3.

Average spacing D between the $I^{\pi} = 1/2^{+}$ levels at $\mathcal{E} \approx B_n$

Nuclei	B_n MeV	D, ev Exper.	D, ev Calc.
^{239}U	4.800	18	31.5
^{169}Er	5.997	125	
^{167}Er	6.436	49	
^{165}Dy	5.715	200	150