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E4-84-117

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IMPROVED DESCRIPTION OF THE FRAGMENTATION OF NUCLEAR COLLECTIVE STATES

Submitted to "ЯФ"

1984

INTRODUCTION

With increasing excitation energy the density of nuclear states increases and their structure becomes complicated. From simple one-quasiparticle or one-phonon states one passes to states with the wave functions containing a lot of components with different number of quasiparticles. One-quasiparticle and one-phonon states are fragmented over many nuclear levels; this fragmentation (distribution of strength) is mainly due to the quasiparticle-phonon interaction. For instance, the widths of giant resonances $\Gamma\downarrow$ are described by taking account of the coupling of 1p-1h and 2p-2h configurations. The quasiparticle-phonon nuclear model [1-3] allows for the wave functions containing one- and two-phonon terms [4-7]. In refs.[8-12] the fragmentation of collective states has been calculated by performing a direct diagonalization in the 2p-2h configuration space, the calculations being restricted by double magic nuclei. The description of the fragmentation of collective states within different models can be found in ref.[13].

It is interesting to consider a possibility of improving the description of the fragmentation of collective one-phonon states. It has been shown in ref.[14] that a more exact inclusion of the Pauli principle leads to small corrections. The one-phonon states forming giant resonances should not be contributed essentially by the effects coming from different types of ground state correlations (see, for instance, ref.[15]). The RPA corrections caused by the terms $a^+a\ a^+a$ turned out to be small[16].

Within the quasiparticle-phonon nuclear model the wave functions describing excited states of doubly even nuclei are represented as a series over the number of phonon operators. In refs.[3-7,14,17,18] only one- and two-phonon terms of the wave functions have been taken into account. General equations of the model taking account of three and four-phonon terms of the wave functions have been obtained in ref.[19]. These equations are very cumbersome. Based on them one can hardly find an approximate method of calculation and to evaluate the influence of the three-phonon terms of the wave functions on the fragmentation of one-phonon states.

In the present paper we give reasons for the necessity of improving the description of the fragmentation of collective states, and the mathematical method is developed for more accurate description of the fragmentation of one-phonon states forming giant resonances.

1. BASIC FORMULAE OF THE MODEL

The model Hamiltonian includes an average nuclear field in the form of a Saxon-Woods potential and the superconducting pairing correlations. It also contains multipole-multipole and spin-multipole - spin-multipole isoscalar and isovector including charge-exchange interactions in the particle-hole and particle-particle channels. In this paper we shall consider only the particle-hole channel. As a basis the model uses the wave functions of one-phonon states calculated in the RPA. In constructing the basis all the parameters are fixed. The solutions of RPA secular equations are taken into account and the model Hamiltonian is transformed as

$$H_M = \sum_{jm} \epsilon_j a^+_{jm} a_{jm} + H_v + H_{vq}, \qquad (1)$$

where $\sum_{jm} \epsilon_j a^+_{jm} a_{jm} + H_v$ describe noninteracting quasiparticles and phonons, H_{vq} is the quasiparticle-phonon interaction. Here ϵ_j is the quasiparticle energy, a^+_{jm} is the quasiparticle creation operator (see ref.[20]), the single-particle states of spherical nuclei are described by quantum numbers jm. The general characteristic of the Hamiltonian is given in ref.[2].

In the calculations performed within the quasiparticle-phonon nuclear model the excited state wave functions of doubly even spherical nuclei are

$$\bar{Q}^+_{JM\nu}\Psi_0, \qquad (2)$$

where Ψ_0 is the phonon vacuum, i.e., $Q_{\lambda\mu i}\Psi_0 = 0$ and also the ground state wave function,

$$\bar{Q}^+_{JM\nu} = \sum_i R_i(J\nu) Q^+_{JMi} + \sum_{\lambda_1 i_1 \lambda_2 i_2} P^{\lambda_1 i_1}_{\lambda_2 i_2}(J\nu)[Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2}]_{JM}, \qquad (3)$$

where $Q^+_{\lambda\mu i}$ is the phonon creation operator and i is the number of the RPA secular equation root. The general form of the system of equations has been obtained in ref.[17]. It was shown that a large set of diagrams was summed. An approximate secular equation for finding the energies $\eta_{J\nu}$ of states (2) disregarding the Pauli principle in the two-phonon terms (3) is

$$\det\left|\left|(\omega_{Ji} - \eta_{J\nu})\delta_{ii'} - \frac{1}{2}\sum_{\lambda_1 i_1 \lambda_2 i_2} \frac{U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)\, U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji')}{\omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2} - \eta_{J\nu}}\right|\right| = 0, \qquad (4)$$

where

$$U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji) = \langle \Psi^*_0 Q_{JMi} H_{vq} [Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2}]_{JM} \Psi_0 \rangle, \qquad (5)$$

$\omega_{\lambda i}$ are the one-phonon state energies calculated in the RPA. In this case the diagrams shown in fig.1a) are summed. Equation (4) serves as a basis for describing the giant resonance widths [3-7,14] and fragmentation of two-quasiparticle states [3] in spherical nuclei.

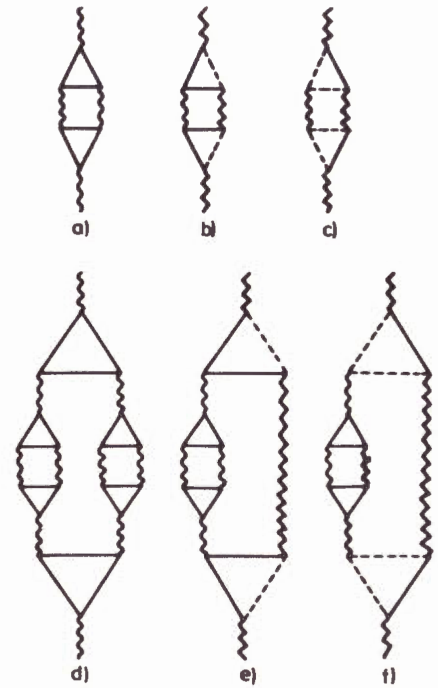

Fig.1. Diagrams in the space of one-phonon states. Notation: ~~~ stands for Q-phonon, ~~~ stands for Ω-phonon, in a) and d) ——— are neutron and photon quasiparticles, in b), c), e) and f) ——— is the neutron quasiparticle, - - - - is the proton quasiparticle.

The fragmentation of neutron-proton np phonons forming charge-exchange resonances is described within the quasiparticle-phonon nuclear model. The wave functions of charge-exchange states are

$$\bar{\Omega}^+_{JM\nu}\Psi_0, \qquad (6)$$

$$\bar{\Omega}_{JM\nu} = \sum_i \bar{R}_i(J\nu)\Omega^+_{JMi} + \sum_{\lambda_1 i_1 \lambda_2 i_2} \bar{P}^{\lambda_1 i_1}_{\lambda_2 i_2}(J\nu)[\Omega^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2}]_{JM}, \qquad (7)$$

$\Omega^+_{\lambda\mu i}$ is the np phonon creation operator determined in ref.[21]. In ref.[18] the general form of the system of basic equations has been obtained and the transition to the system of approximate equations has been made. The approximate secular equation for finding the energies of states (6) disregarding the Pauli principle in the two-phonon terms (7) is

$$\det\|(\Omega_{Ji}-\zeta_{J\nu})\delta_{ii'} - \sum_{\lambda_1 i_1 \lambda_2 i_2} \frac{V^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)\, V^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji')}{\Omega_{\lambda_1 i_1}+\omega_{\lambda_2 i_2}-\zeta_{J\nu}}\| = 0, \qquad (8)$$

where

$$V^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji) = \langle \Psi^*_0\, \Omega_{JMi} H_{vq} [\Omega^+_{\lambda_1 \mu_1 i_1} Q^+_{\lambda_2 \mu_2 i_2}]_{JM} \Psi_0 \rangle, \qquad (9)$$

and Ω_{λ_i} are the energies of neutron-proton phonons. In this case the diagrams from fig.1b) and c) are summed. Equation (8) is used for describing the fragmentation of charge-exchange resonances in spherical nuclei /22/.

2. IMPROVED DESCRIPTION OF THE FRAGMENTATION OF CHARGE-EXCHANGE ONE-PHONON STATES

Within the quasiparticle-phonon nuclear model an excited state wave function is represented as a series over the number of phonons. It is extremely difficult to investigate the influence of many-photon terms of the wave functions on the fragmentation of one-phonon states. This problem is of the same difficulty as the elucidation of the role of rejected within the Hartree-Fock-Bogolubov approximation chains of equations of the many-body problem. In both the cases the approximate equations are asserted to describe correctly the basic characteristics of nuclear excitations and the rejected terms are partially taken into account by using some constants fixed from the experimental data. Undoubtedly one should make attempts at calculating corrections due to rejected numbers.

The influence of the quasiparticle plus two phonons terms of the wave functions on the fragmentation of one-quasiparticle states has been studied in ref. /23/. It was shown that the inclusion of these terms improves the description of deep hole states in spherical nuclei. The calculation of the radiative strength functions /24/ with the wave function containing one-quasiparticle, quasiparticle plus phonon and quasiparticle plus two phonons terms shows that the fragmentation of one-phonon states due to the coupling with two-phonon ones through the functions $U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)$ (see formula (5)) changes strongly the reduced E1 and M1 transition probabilities. Such a change is exemplified in fig.2. It is seen from this figure that the fragmentation of 1^+ one-phonon states causes a strong change in the behaviour of the M1 strength function describing the transitions from neutron resonances to the ground state of ^{61}Ni. It has been shown in refs. /6,7/ that due to the fragmentation of

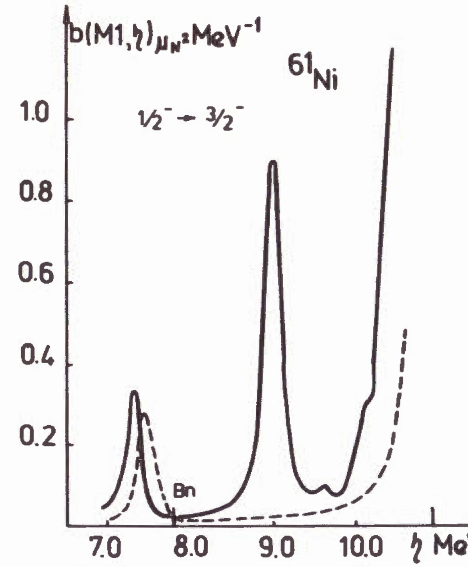

Fig.2. Strength functions of the M1 transitions in ^{61}Ni from p-wave resonances with $J^\pi = 1/2^-$ to the ground state with $J^\pi = 3/2^-$. Notation: solid line is the calculations allowing for the function $U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)$ (formula (5)), dashed line is the calculations disregarding $U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)$.

one-phonon 1^- states there appear substructures in the photoabsorption cross sections on spherical nuclei on the giant dipole resonance tails and part of the giant dipole resonance strength is pushed into the region of (8-10) MeV.

These examples show that the fragmentation of one-phonon states turns out to be essential in the excitation energy region of 8-10 MeV. One should remember that the quasiparticle-phonon interaction couples most strongly the terms of the wave functions which differ by one phonon. Therefore, the inclusion of three-phonon terms of the wave functions will influence the fragmentation of two-phonon terms, whereas the latter will enhance the fragmentation of one-phonon states. There arises a question whether the influence of the fragmentation of two-phonon terms on that of one-phonon terms can be taken into account without calculating directly the coupling with the three-phonon terms of the wave functions.

The main point of an improved description of the fragmentation of one-phonon states is the use of the one-phonon states, which have already been fragmented, in the two-phonon terms of the wave functions. In this case one does not pass to new phonon operators. As earlier, the RPA phonons and the corresponding vacuum are used as a basis. This method is aimed at effective account of some three-phonon terms of the wave functions, i.e., restricted number of 3p-3h states.

The study of the fragmentation of the GT resonance in such spherical nuclei as ^{90}Zr and ^{208}Pb demonstrates a more accurate description. The inclusion of fragmented phonons in the two-phonon part of the wave function may appear to be essential. First, we shall get formulae for a more accurate description of the fragmentation of charge-exchange one-phonon states.

The wave function of an odd-odd nucleus is

$$\Psi_n(JM) = \{ \sum_i \bar{C}_i(Jn)\Omega^+_{JMi} + \sum_{J_1 i_1 J_2 \nu_2} \bar{D}^{J_1 i_1}_{J_2 \nu_2}(Jn)[\Omega^+_{J_1 M_1 i_1} \bar{Q}^+_{J_2 M_2 \nu_2}]_{JM} \}\Psi_0 , \quad (10)$$

where the operator $\bar{Q}^+_{JM\nu}$ is determined by formula (3). Let eqs. (4) be solved and the functions entering these operators be determined. The normalization condition (10) disregarding the Pauli principle in many-phonon components is

$$\sum_i (\bar{C}_i(Jn))^2 + \sum_{J_1 i_1 J_2 \nu_2} (\bar{D}^{J_1 i_1}_{J_2 \nu_2}(Jn))^2 = 1 . \quad (11)$$

Under such a choice of the wave function (10) certain three-phonon terms are taken into account. An average value of H_M over (10) is

$$<\Psi^*_n(JM) H_M \Psi_n(JM)> = \sum_i \Omega_{Ji}(\bar{C}_i(Jn))^2 + \sum_{J_1 i_1 J_2 \nu_2} (\Omega_{J_1 i_1} + \eta_{J_2 \nu_2}) \cdot$$
$$(\bar{D}^{J_1 i_1}_{J_2 i_2}(Jn))^2 + 2 \sum_{i J_1 i_1 J_2 \nu_2} \bar{C}_i(Jn) \bar{D}^{J_1 i_1}_{J_2 \nu_2}(Jn) \tilde{V}^{J_1 i_1}_{J_2 \nu_2}(Ji) , \quad (12)$$

$$\tilde{V}^{J_1 i_1}_{J_2 \nu_2}(Ji) = \sum_{i_2} R_{i_2}(J_2 \nu_2) V^{J_1 i_1}_{J_2 i_2}(Ji) . \quad (13)$$

Using the variational principle we get the following equations:

$$(\Omega_{Ji} - E_n)\bar{C}_i(Jn) + \sum_{J_1 i_1 J_2 \nu_2} \tilde{V}^{J_1 i_1}_{J_2 \nu_2}(Ji) \bar{D}^{J_1 i_1}_{J_2 \nu_2}(Jn) = 0 , \quad (14)$$

$$(\Omega_{J_1 i_1} + \eta_{J_2 \nu_2} - E_n)\bar{D}^{J_1 i_1}_{J_2 \nu_2}(Jn) + \sum_{i'} \tilde{V}^{J_1 i_1}_{J_2 \nu_2}(Ji')\bar{C}_{i'}(Jn) = 0 . \quad (14')$$

Equations (11), (14) and (14´) form a complete system for finding the functions $\bar{C}, \bar{D}$ and energies E_n of the states described by the wave functions (10). Substituting $\bar{D}$ from (14´) into (14), one gets the following secular equation:

$$\mathfrak{F}^{CH}_J(E_n) \equiv \det \| (\Omega_{Ji} - E_n)\delta_{ii'} - \sum_{J_1 i_1 J_2 \nu_2} \frac{\tilde{V}^{J_1 i_1}_{J_2 \nu_2}(Ji) \tilde{V}^{J_1 i_1}_{J_2 \nu_2}(Ji')}{\Omega_{J_1 i_1} + \eta_{J_2 \nu_2} - E_n} \| = 0 . \quad (15)$$

The rank of this determinant is equal to the number of one-phonon terms in the wave function (10). In this case the diagrams of the type of fig.1e) and f) are summed. In the numerical calculations part of fragmented phonons can be changed by one-phonon states.

Charge-exchange collective states are excited in reactions of the type (p,n) and (n,p). The amplitudes of the relevant transitions are expressed through $\Phi^{(\mp)}_{\lambda i}$ the explicit form of which is given in ref. /22/. The calculations within the QPM are based on the strength function by which the strength distribution of certain physical quantities averaged over the energy interval Δ is calculated without solving eq. (15). The strength functions for the calculation of strength distribution of the (p,n) and (n,p) transitions in the approximation under consideration are

$$b^{(\mp)}(\Phi, E) = \frac{1}{\pi} \mathrm{Im} \{ \frac{\sum_{ii'} \tilde{A}^{CH}_{ii'}(E + i\frac{\Delta}{2})\Phi^{(\mp)}_{Ji}\Phi^{(\mp)}_{Ji'}}{\mathfrak{F}^{CH}_J(E + i\frac{\Delta}{2})} \} , \quad (16)$$

$\tilde{A}^{CH}_{ii'}$ is the minor of determinant (15).

As a result of solution of eq.(4) for the states with fixed values of J^π we find the energies $\eta_{J\nu}$. The number of states is equal to the number of one- and two-phonon poles. The strength of each one-phonon state i is distributed over several levels ν. So, in ^{118}Sn in the energy interval of 7-17 MeV there are two one-phonon states with $J^\pi = 2^+$ with large B(E2)-values. The fragmentation of these one-phonon states with energies $\omega_{2i=1}$ = 11.97 MeV and $\omega_{2i=2}$ =12.33 MeV over the levels ν in ^{118}Sn is shown in fig.3. Instead of two poles in eq. $\Omega_{J_1 i_1} + \omega_{21}$ and $\Omega_{J_1 i_1} + \omega_{22}$, in the secular equation (15) one should take into account about 200 poles $\Omega_{J_1 i_1} + \eta_{J_2 \nu_2}$. It is seen from the figure that there are solutions in the energy intervals 10.4-10.8 MeV, about 11 MeV, 11.2-11.6 MeV, 11.8-12.9 MeV and so on. One can take only a part of these equations into account, for instance, 60 levels over which more than 90% of strength of these one-phonon states is distributed. For each fixed state J = 2, $\nu = \nu_0$ with energy $\eta_{2\nu_0}$ there are calculated functions $R_i(2\nu_0)$ entering in (13). In the case under consideration the summation is over two states i = 1 and 2. In the rest terms in (15) the summation over $J_2 \nu_2$ is changed by $J_2 i_2$ as in (4).

The effect of the inclusion of fragmented one-phonon states, as is shown in fig.3, may be essential since the number of levels over which the fragmentation proceeds is very large and their energy interval is broad. The strength of charge-exchange resonances at their maxima may be reduced by taking into account this effect.

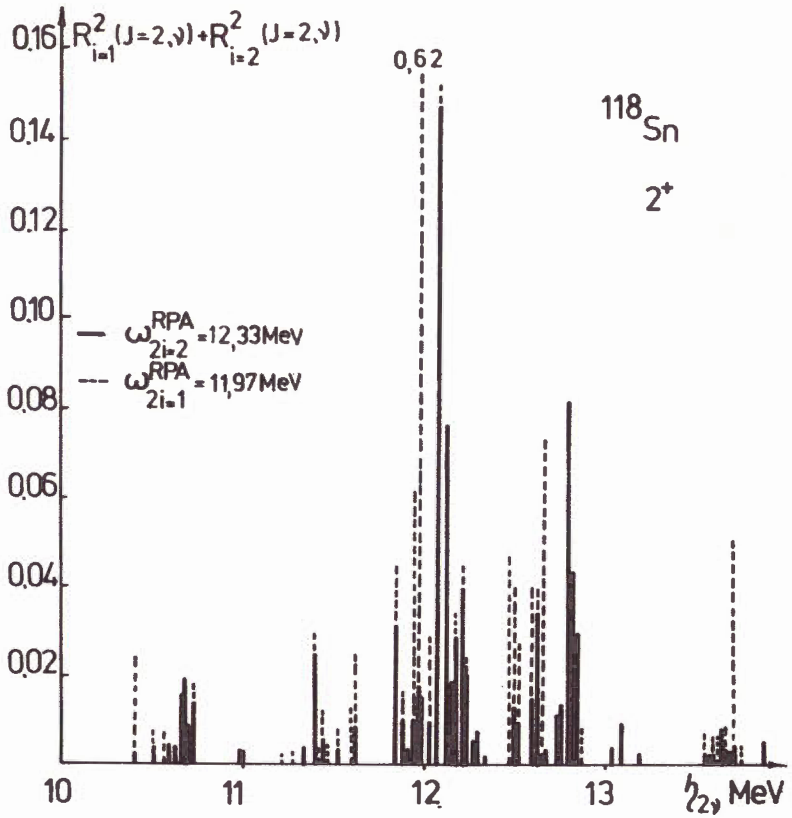

Fig.3. Distribution of strength of the one-phonon $R^2_{i=1}(2\nu)$ with energy $\omega_{2i=1} = 11.97$ MeV (denoted by ---) and $R^2_{i=2}(2\nu)$ with $\omega_{2i=2} = 12.33$ MeV (denoted by ——) states with $J^\pi = 2^+$ in ^{118}Sn over the solutions of the secular equation (4).

3. ANOTHER VERSION OF IMPROVING THE DESCRIPTION OF THE FRAGMENTATION OF ONE-PHONON STATES

Let us consider the fragmentation of a one-phonon state in such an approximation when both phonons in the two-phonon part of the wave function are fragmented. In this case we take into account a part of three- and four-phonon terms in the wave function. The excited state wave function of a doubly even spherical nucleus is

$$\Psi_n(JM) = \{ \sum_i C_i(Jn) Q^+_{JMi} + \sum_{J_1\nu_1 J_2\nu_2} D^{J_1\nu_1}_{J_2\nu_2}(Jn) [\tilde{Q}^+_{J_1M_1\nu_1} \tilde{Q}^+_{J_2M_2\nu_2}]_{JM} \}\Psi_0 . \quad (17)$$

Consider that the secular eq.(4) is solved and the functions $R_i(J\nu)$ and energies $\eta_{J\nu}$ entering into operator $\tilde{Q}^+_{JM\nu}$ are found. Its normalization condition without the Pauli principle in many-phonon components is

$$\sum_i C_i^2(Jn) + 2 \sum_{J_1\nu_1 J_2\nu_2} (D^{J_1\nu_1}_{J_2\nu_2}(Jn))^2 = 1 . \quad (18)$$

An average value of H_M over (17) is

$$<\Psi^*_n(JM) H_M \Psi_n(JM)> = \sum_i \omega_{Ji} C_i^2(Jn) + 2 \sum_{J_1\nu_1 J_2\nu_2} (\eta_{J_1\nu_1} + \eta_{J_2\nu_2}) \cdot$$

$$(D^{J_1\nu_1}_{J_2\nu_2}(Jn))^2 + 2 \sum_i \sum_{J_1\nu_1 J_2\nu_2} C_i(Jn) D^{J_1\nu_1}_{J_2\nu_2}(Jn) \tilde{U}^{J_1\nu_1}_{J_2\nu_2}(Ji).$$

Using the variational principle we get the secular equations

$$(\omega_{Ji} - E_n) C_i(Jn) + \sum_{J_1\nu_1 J_2\nu_2} D^{J_1\nu_1}_{J_2\nu_2}(Jn) \tilde{U}^{J_1\nu_1}_{J_2\nu_2}(Ji) = 0 , \quad (19)$$

$$(\eta_{J_1\nu_1} + \eta_{J_2\nu_2} - E_n) D^{J_1\nu_1}_{J_2\nu_2}(Jn) + \frac{1}{2} \sum_{i'} C_{i'}(Jn) \tilde{U}^{J_1\nu_1}_{J_2\nu_2}(Ji') = 0 . \quad (20)$$

Equations (18), (19) and (20) form a complete system of equations for finding the functions C , D and energies E_n of the states described by the wave functions (17). The secular equation is

$$\tilde{\mathcal{F}}_J(E_n) \equiv \det \| (\omega_{Ji} - E_n)\delta_{ii'} - \frac{1}{2} \sum_{J_1\nu_1 J_2\nu_2} \frac{\tilde{U}^{J_1\nu_1}_{J_2\nu_2}(Ji) \tilde{U}^{J_1\nu_1}_{J_2\nu_2}(Ji')}{\eta_{J_1\nu_1} + \eta_{J_2\nu_2} - E_n} \| = 0, \quad (21)$$

where

$$\tilde{U}^{J_1\nu_1}_{J_2\nu_2}(Ji) = \sum_{i_1 i_2} R_{i_1}(J_1\nu_1) R_{i_2}(J_2\nu_2) U^{J_1 i_1}_{J_2 i_2}(Ji) . \quad (22)$$

The rank of the determinant (21) coincides with that of the determinant (4) and equals the number of one-phonon states in the wave functions (2) or (17). The energies $\eta_{J\nu}$ and functions $R_i(J\nu)$ calculated earlier, which determine the fragmentation of one-phonon states, are used to calculate (21) and (22). In this case the diagrams of the type of fig.1d) are summed. The calculations can be performed also in the approximation when part of the states in the sum over $J_1\nu_1$ and $J_2\nu_2$ is substituted by the one-phonon states $J_1 i_1$ and $J_2 i_2$ for which one of the functions R is equal to unity and the others to zero and the energies $\eta_{J_1\nu_1}$ and $\eta_{J_2\nu_2}$ are substituted by $\omega_{J_1 i_1}$ and $\omega_{J_2 i_2}$. Thus, in the numerical calculations some of the phonons in the two-phonon part of the wave function (17) may be already fragmented and the others nonfragmented. This means that calculations can be performed with the wave function

$$\Psi_n(JM) = \{ \sum_i C_i(Jn) Q^+_{JMi} + \sum_{J_1\nu_1 J_2\nu_2} D^{J_1\nu_1}_{J_2\nu_2}(Jn) [\tilde{Q}^+_{J_1M_1\nu_1} \tilde{Q}^+_{J_2M_2\nu_2}]_{JM} + \sum_{\lambda_1 i_1 \lambda_2 i_2} D^{\lambda_1 i_1}_{\lambda_2 i_2}(Jn) [Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2}]_{JM} \} \Psi_0 , \quad (23)$$

where sets of states $J_1\nu_1 J_2\nu_2$ and $\lambda_1 i_1 \lambda_2 i_2$ are chosen so as to avoid double counting.

Now we find the strength function describing excitation of states with the wave functions (17) from the ground state of the relevant nucleus. Let excitation of the state $\Psi_n(JM)$ proceed from the ground state through one-phonon components with amplitude

$$\Phi_{Jn} = \sum_i C_i(Jn) \Phi_{Ji} . \quad (24)$$

A form of Φ_{Ji} depends on the process under consideration. For instance, in the excitation of the state $\Psi_n(JM)$ by γ-quanta Φ_{Ji} is determined by the operators of the $E\lambda$ - or $M\lambda$ -transitions (see ref.[8]) from the ground to one-phonon states. The corresponding strength function is

$$b(\Phi, E) = \frac{1}{\pi} \operatorname{Im} \{ \frac{\sum_{ii'} \tilde{A}_{ii'}(E + i\frac{\Delta}{2}) \Phi_{Ji} \Phi_{Ji'}}{\tilde{\mathcal{F}}_J (E + i\frac{\Delta}{2})} \} , \quad (25)$$

where $\tilde{A}_{ii'}$ is the minor of the determinant (21). The strength function(25)coincides in form with the strength function for(2). The difference is in the functions $\tilde{\mathcal{F}}(E + i\frac{\Delta}{2})$ and $\tilde{A}_{ii'}(E + i\frac{\Delta}{2})$.

An approximate method developed in this paper allows one to take into account a part of the influence of the 3p-3h configurations on the fragmentation of one-phonon states in doubly even spherical nuclei. The numerical calculations by the formulae given above can be made in the nearest future at the following generation of computers.

The author is grateful to A.I.Vdovin, V.V.Voronov, L.A.Malov and I.N.Mikhailov for interesting discussions and to V.Yu.Ponomarev for help in making calculations.

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Received by Publishing Department on February 23,1984.

Соловьев В.Г. E4-84-117
Улучшение описания фрагментации ядерных коллективных состояний

Разработан математический метод более точного описания фрагментации однофононных состояний, формирующих гигантские резонансы. Идея метода заключается в использовании в двухфононных членах волновых функций уже фрагментированных однофононных состояний. Получены силовые функции для возбуждения коллективных зарядово-обменных состояний и гигантских резонансов в сферических ядрах.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна 1984

Soloviev V.G. E4-84-117
Improved Description of the Fragmentation of Nuclear Collective States

A mathematical method is developed for a more accurate description of the fragmentation of one-phonon states forming giant resonances. The method consists in that the one-phonon states already fragmented are used in the two-phonon terms of wave functions. Strength functions are obtained for the excitation of collective charge-exchange states and giant resonances in spherical nuclei.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

Preprint of the Joint Institute for Nuclear Research. Dubna 1984