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Соловьев В.Г., Стоянов Ч., Вдовин А.И. **E4 - 10397**

Изучение фрагментации гигантских мультипольных резонансов по двухфононным состояниям в сферических ядрах

Изложен метод вычисления силовых функций фотовозбуждения гигантских мультипольных резонансов в сферических ядрах в рамках модели, основанной на взаимодействии квазичастиц с фононами. Однофононные состояния вычислены в приближении хаотичных фаз. Рассчитана фрагментация однофононных состояний по двухфононным. Получены силовые функции $b(E\lambda,\eta)$, характеризующие энергетические области нахождения ГМР в ядрах ^{90}Zr, ^{120}Sn и ^{124}Te. Изучались изовекторный дипольный, изоскалярный и изовекторный квадрупольные и низкоэнергетический октупольный резонансы.

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Soloviev V.G., Stoyanov Ch., Vdovin A.I. **E4 - 10397**

Fragmentation of Giant Multipole Resonances Over Two-Phonon States in Spherical Nuclei

A method for calculation of the strength function photoexcitation of giant multipole resonances in spherical nuclei is presented. The method is developed in the framework of the model based on the quasiparticle phonon interaction. The one-phonon states are calculated in the random phase approximation. The fragmentation of one-phonon states over two-phonon ones is calculated. The strength functions $b(E\lambda,\eta)$ and the positions of giant isovector dipole and isoscalar quadrupole resonances in ^{90}Zr, ^{120}Sn and ^{124}Te are found. The calculations of the transition strengths and positions of these resonances are in reasonable agreement with experiment. The strength function $b(E2,\eta)$ is calculated for the isovector quadrupole resonance in ^{90}Zr. The description of the low-energy octupole resonances in ^{90}Zr and ^{120}Sn is in agreement with experiment.

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1. INTRODUCTION

Most progress has recently been achieved in the experimental study of giant multipole resonances. In many nuclei the new resonances, isoscalar E2 and M1 resonance, are observed. There is some evidence for the existence of an isovector E2 strength and a low-energy octupole resonance [1-5]. The experimental success has stimulated theoretical studies in these lines. A remarkable feature of relevant theoretical studies [6-10] is the application of the semi-microscopic methods earlier used for the description of low-lying nuclear excitations. Many calculations were performed within the microscopic method using the effective nucleon-nucleon forces [11,12]. As a rule, the aim of these calculations was to obtain the resonance positions, the total photoexcitation probability, and the resonance contributions to the energy weighted sum rules (EWSR).

For spherical nuclei in the random phase approximation (RPA) the giant resonances are formed by one or a few one-phonon states spaced by several hundreds of keV. But for deformed nuclei the situation is quite different [8]. Here the resonances are formed by many (several dozens or even hundreds) one-phonon states within the region almost coinciding with the experimental one. The large width of the spherical nucleus photoabsorption spectra is mainly due to the fragmentation of collective one-phonon states (or particle-hole states) over more complex configurations. The fragmentation is caused by coupling of different modes of nucleus excitations. It was pointed out earlier[13-15] that the spreading of the giant resonance strength is connected with the low energy collective properties of the nucleus. Besides, due to high excitation energy, the decay channels with one or a few outgoing nucleons are very important. Therefore, the continuous spectrum should be taken into account [11,12]. The description of the giant resonance width within the semi-microscopic and microscopic models is, for the time being, at the initial stage [12,16,17].

The aim of this paper is to calculate the fragmentation of the one-phonon states, forming a giant resonance, within the RPA over two-phonon states. The strength functions $b(E\lambda,\eta)$ of the $E\lambda$ resonance excitation are calculated as functions of the excitation energy η. The isovector dipole, isoscalar and isovector quadrupole and low-energy octupole resonances are considered in nuclei ^{90}Zr, ^{120}Sn, and ^{124}Te.

2. THE MODEL

For our calculation we use a modified version of the semi-microscopic nuclear model[18,19]. This model was used for describing the level density[20] and the neutron strength functions in medium and heavy nuclei[21-23]. It gives also satisfactory results for the giant resonances in deformed nuclei in RPA [8,23].

Let us discuss the main assumptions of the model. The model Hamiltonian includes the average field for protons and neutrons, superconducting pairing interactions and multipole-multipole and spin-multipole-spin-multipole forces. Both the isovector and isoscalar multipole-multipole forces have been used. The isovector forces are used for obtaining the isovector dipole and quadrupole resonances. The isovector forces change noticeably the distributions of $B(E\lambda)$ values at high excitation energies (10-15 MeV)[8,23] and, in some cases, at intermediate excitation energies.

We write the Hamiltonian in terms of the quasiparticle ($\alpha^+_{jm}, \alpha_{jm}$) and phonon ($Q^+_{\lambda\mu i}, Q_{\lambda\mu i}$) operators as follows

$$H \equiv H_{ph} + H_{qph} = \sum_{\lambda\mu i} \omega_{\lambda i} Q^+_{\lambda\mu i} Q_{\lambda\mu i} - \frac{1}{2\sqrt{2}} \sum_{\lambda\mu i} (-)^{\lambda-\mu} [Q^+_{\lambda\mu i} \pm (-)^{\lambda-\mu} Q_{\lambda-\mu i}] \times \tag{1}$$

$$\times \Big\{ \sum^{n}_{j_1 j_2} \frac{f^{\lambda}_{j_1 j_2} v^{(\mp)}_{j_1 j_2}}{\sqrt{\mathcal{Y}_n(\lambda i)}} B(j_1 j_2 \lambda -\mu) + \sum^{p}_{j_1 j_2} \frac{f^{\lambda}_{j_1 j_2} v^{(\mp)}_{j_1 j_2}}{\sqrt{\mathcal{Y}_p(\lambda i)}} B(j_1 j_2 \lambda -\mu) \Big\} +$$

$$+ \text{h.c.},$$

where

$$B(j_1 j_2 \lambda\mu) = \sum_{m_1 m_2} \langle j_1 m_1 j_2 m_2 | \lambda\mu \rangle (-)^{j_2+m_2} \alpha^+_{j_1 m_1} \alpha_{j_2 -m_2} .$$

We use the following notation: $f^{\lambda}_{j_1 j_2}$ is the reduced matrix element of the multipole (or spin-multipole) operator:

$$u^{(\pm)}_{j_1 j_2} = u_{j_1} v_{j_2} \pm u_{j_2} v_{j_1}, \quad v^{(\pm)}_{j_1 j_2} = u_{j_1} u_{j_2} \pm v_{j_1} v_{j_2},$$

u_j, v_j are the Bogolubov transformation coefficients. In eq. (1) and in what follows the indices n(p) mean that the summation runs over the neutron (proton) single-particle states. The energy $\omega_{\lambda i}$ of the one-phonon state $Q^+_{\lambda\mu i}|0\rangle_{ph}$ is found by solving the following equation[18]:

$$\frac{\kappa_0^\lambda + \kappa_1^\lambda}{2\lambda+1}(X_n^{\lambda i} + X_p^{\lambda i}) - \frac{4\kappa_0^\lambda \kappa_1^\lambda}{(2\lambda+1)^2} X_n^{\lambda i} X_p^{\lambda i} - 1 = 0, \qquad (2)$$

where κ_0^λ, κ_1^λ are the isoscalar and isovector constants of the multipole-multipole interactions, and

$$X^{\lambda i}_{n,p} = \sum_{j_1 j_2}^{n,p} \frac{[f^\lambda_{j_1 j_2} u^{(\pm)}_{j_1 j_2}]^2 \epsilon_{j_1 j_2}}{\epsilon^2_{j_1 j_2} - \omega^2_{\lambda i}}, \qquad (3)$$

$\epsilon_{j_1 j_2}$ is the two-quasiparticle state energy. In eqs. (1) and (3) $u^{(+)}_{j_1 j_2}$ and $v^{(-)}_{j_1 j_2}$ correspond to the multipole phonons and $u^{(-)}_{j_1 j_2}$ and $v^{(+)}_{j_1 j_2}$ to the spin-multipole ones. Determining the quantities $Y_{n,p}(\lambda i)$ from the following relations:

$$Y_{n,p}(\lambda i) = \frac{1}{2} \frac{\partial}{\partial \omega} X^{\lambda i}_{n,p} \Big|_{\omega = \omega_{\lambda i}}$$

for $\mathcal{Y}_{n,p}(\lambda i)$ we have

$$\mathcal{Y}_{n,p}(\lambda i) = Y_{n,p}(\lambda i) + Y_{p,n}(\lambda i)\left\{\frac{1 - \frac{\kappa_0^\lambda + \kappa_1^\lambda}{2\lambda+1} X^{\lambda i}_{n,p}}{\frac{\kappa_0^\lambda - \kappa_1^\lambda}{2\lambda+1} X^{\lambda i}_{p,n}}\right\}^2. \qquad (4)$$

In eq. (1) the term H_{qph} describes the quasiparticle-phonon interaction, i.e., in our model it favours the fragmentation of one-phonon states over more complex ones. It is clearly seen from eqs. (1)-(4) that the quasiparticle-phonon interaction does not contain any new parameters.

We take into account the coupling between the one-phonon and two-phonon states. Thus, the wave function of a highly excited state is

$$\Psi_\nu(JM) = \{\sum_i R_\nu(Ji) Q^+_{JMi} + \sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} P^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)[Q^+_{\lambda_1 \mu_1 i_1} Q^+_{\lambda_2 \mu_2 i_2}]_{JM}\}|0\rangle_{ph}, \qquad (5)$$

where $|0\rangle_{ph}$ is the phonon vacuum and, at the same time, the ground state of a doubly even nucleus. It should be noted that the two-phonon state density exceeds that of the two particles-two holes (2p-2h) states.

The secular equation has the form:

$$\mathcal{F}(\eta) \equiv \det|(\omega_{Ji} - \eta_{J\nu})\delta_{ii'} - \sum_{\substack{\lambda_1 i_1\\ \lambda_2 i_2}} \frac{U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji) U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji')}{\omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2} - \eta_{J\nu}}| = 0. \qquad (6)$$

Its roots are the energies, $\eta_{J\nu}$, of the states $\Psi_\nu(JM)$ (see refs.[18,24]). The matrix rank in eq. (6) coincides with the number of the one-phonon states in the one-phonon part of the wave function (5). The coefficients $U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji)$ are of the following form:

$$U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji) = \langle 0|Q_{JMi} H_{qph}[Q^+_{\lambda_1 \mu_1 i_1} Q^+_{\lambda_2 \mu_2 i_2}]_{JM}|0\rangle \equiv$$

$$\equiv U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji, n) + U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji, p),$$

where, for instance,

$$U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji,n) = \frac{1}{\sqrt{2}}[(2\lambda_1+1)(2\lambda_2+1)]^{1/2}(-)^{\lambda_1+\lambda_2+J} \times$$

$$\times \sum_{j_1 j_2 j_3} \frac{f^{J}_{j_1 j_2} v^{(\mp)}_{j_1 j_2}}{\sqrt{\mathcal{Y}_n(Ji)}} \begin{Bmatrix} \lambda_1 & \lambda_2 & J \\ j_2 & j_1 & j_3 \end{Bmatrix} (\phi^{\lambda_1 i_1}_{j_3 j_1} \psi^{\lambda_2 i_2}_{j_2 j_3} + \phi^{\lambda_2 i_2}_{j_2 j_3} \psi^{\lambda_1 i_1}_{j_3 j_1}) +$$

$$+ \frac{f^{\lambda_1}_{j_1 j_2} v^{(\mp)}_{j_1 j_2}}{\sqrt{\mathcal{Y}_n(\lambda_1 i_1)}} \begin{Bmatrix} \lambda_1 & \lambda_2 & J \\ j_3 & j_2 & j_1 \end{Bmatrix} (\psi^{\lambda_2 i_2}_{j_3 j_1} \psi^{Ji}_{j_2 j_3} + \phi^{\lambda_2 i_2}_{j_3 j_1} \phi^{Ji}_{j_2 j_3}) +$$

$$+ \frac{f^{\lambda_2}_{j_1 j_2} v^{(\mp)}_{j_1 j_2}}{\sqrt{\mathcal{Y}_n(\lambda_2 i_2)}} \begin{Bmatrix} \lambda_1 & \lambda_2 & J \\ j_1 & j_3 & j_2 \end{Bmatrix} (\psi^{Ji}_{j_3 j_1} \psi^{\lambda_1 i_1}_{j_2 j_3} + \phi^{Ji}_{j_3 j_1} \phi^{\lambda_1 i_1}_{j_2 j_3}) \quad (7)$$

and the formula for $U^{\lambda_2 i_2}_{\lambda_1 i_1}(Ji,p)$ follows from eq. (7) by the substitution $n \to p$. The functions $\psi^{\lambda i}_{j_1 j_2}$ and $\phi^{\lambda i}_{j_1 j_2}$ are the forthgoing and back-going amplitudes of one-phonon states. As far as we are dealing with the fragmentation of one-phonon states only, we write the expressions for $R_\nu(Ji)$ as

$$R^2_\nu(Ji) = \frac{M^2_{ii}}{\sum_k (M_{ik})^2 + \frac{1}{2} \sum_{\substack{\lambda_1 i_1 \\ \lambda_2 i_2}} \frac{(\sum_k M_{ik} U^{\lambda_2 i_2}_{\lambda_1 i_1}(Jk))^2}{\omega_{\lambda_1 i_1} + \omega_{\lambda_2 i_2} - \eta_{J\nu}}}, \quad (8)$$

M_{ik} is the minor of the determinant in eq. (6). Formula (8) can be rewritten as

$$R^2_\nu(Ji) = -\left(\frac{\partial}{\partial \eta} \frac{\mathcal{F}(\eta)}{M_{ii}}\right)^{-1}\Big|_{\eta = \eta_{J\nu}}. \quad (9)$$

It is very difficult to obtain numerous solutions of eq. (6) when many one-phonon terms in the wave function (5) are taken into account. On the other hand, we do not need such a large amount of information obtained by solving eq. (6) and calculating the structure of each of many hundreds of states. Therefore it seems to us more reasonable to calculate the averaged resonance characteristics observed, like the photoexcitation probability, without solving eq. (6).

To calculate the strength function $b(E\lambda,\eta)$ we apply the method which was earlier used by Bohr and Mottelson[25] and the authors of refs.[8,21,22]. We introduce the function

$$b(EJ,\eta) = \sum_\nu \rho(\eta - \eta_{J\nu}) B(EJ, 0^+_{g.s.} \to J_\nu),$$

were the sum is taken over all the roots of eq. (6) in the excitation energy interval ΔE, and

$$\rho(\eta - \eta_{J\nu}) = \frac{1}{2\pi} \frac{\Delta}{(\eta-\eta_\nu)^2 + \frac{\Delta^2}{4}}.$$

The function $b(EJ,\eta)$ can be written as a contour integral around the poles which are the roots of eq. (6). After some transformations[25,26] we get

$$b(EJ,\eta) = \frac{e^2}{2\pi} \mathrm{Im} \frac{\Phi^2(\eta + i\frac{\Delta}{2})}{\mathcal{F}(\eta + i\frac{\Delta}{2})}, \quad (10)$$

$$\Phi^2(\eta) = \sum_{ik} M_{ik} \Phi_i \Phi_k, \quad (11)$$

$$\Phi_i = e^{(J)} \frac{X^{Ji}_n}{\sqrt{\mathcal{Y}_n(Ji)}} + (1 + e^{(J)}) \frac{X^{Ji}_p}{\sqrt{\mathcal{Y}_p(Ji)}}, \quad (12)$$

where $e^{(J)}$ is the effective charge. It should be mentioned that, when deriving eq. (10), we neglect the terms $\sim \alpha^+\alpha$ in the $E\lambda$ transition operator.

Using eq. (10) we can immediately calculate $b(EJ,\eta)$ as a function of the excitation energy without solving eq. (6) and calculating the structure of each state. The problem becomes simpler when a new parameter Δ is introduced. If it is small enough $\Delta \ll D$, where D is the average spacing of the two-phonon states, the curve $b(EJ,\eta)$ exhibits as a series of narrow peaks of the Breit-Wigner form, the center of each coinciding with the location of the appropriate root of eq. (6). With increasing Δ the peaks become broader, their amplitudes decrease and the neighbouring peaks overlap. The fine structure gradually disappears. Thus, by varying Δ, we can reproduce the B(EJ) distribution over the roots of eq. (6) in as much detail as we wish. We need a picture of the B(EJ) distribution which conserves only the main features of the solution of eq. (6). This picture remains almost unaffected with varying Δ in the interval 200-500 keV.

We use energy weighted sum rules (EWSR). The model-independent dipole EWSR is

$$\sum_\nu b(E1,\eta_\nu)\eta_\nu = \frac{9}{8\pi}\frac{e^2\hbar^2}{m}\frac{NZ}{A} = 0.18\frac{N \cdot Z}{A}e^2 b \text{ MeV}. \quad (13)$$

Since both the isoscalar and isovector components for the resonances with $\lambda > 1$ are simultaneously taken into account, we use the sum rule in the form

$$\sum_\nu b(E\lambda,\eta_\nu)\eta_\nu = 1.65\lambda(2\lambda+1)^2 Z \langle r^{2\lambda-2}\rangle e^2 \text{fm}^{2\lambda} \text{MeV} =$$
$$= 4.95\lambda(2\lambda+1)(1.2)^{2\lambda-2} Z A^{\frac{2\lambda-2}{3}} 10^{-2\lambda} e^2 b^\lambda \text{ MeV}. \quad (14)$$

In eq. (14) $\langle r^{2\lambda-2}\rangle$ is evaluated, as usual, for the square well with radius $1.2A^{1/3}$ fm.

Separating the isoscalar part of eq. (14), we have

$$\sum_\nu b(E\lambda,\eta_\nu)\eta_\nu = 4.95\lambda(2\lambda+1)(1.2)^{2\lambda-2} Z^2 A^{\frac{2\lambda-5}{3}} 10^{-2\lambda} e^2 b^\lambda \text{MeV}. \quad (15)$$

3. NUMERICAL DETAILS

In order to describe correctly the structure of giant resonances at an excitation energy of several dozens of MeV, it is necessary to consider a large number of single-particle levels and, in particular, the quasicontinuous spectrum levels. We use the Saxon-Woods potential [27]. Its parameters are given in Table 1. The single-particle energies and the matrix elements of multipole and spin-multipole operators are calculated with the aid of the programme suggested in ref. [28] and composed on the basis of the method of ref [29]. Figure 1

Table 1
Single-particle Saxon-Woods potential parameters

		r_0 fm	V_0 MeV	κ fm^2	a fm^{-1}
A=91	N	1.29	45.0	0.413	1.613
Z=39	Z	1.24	57.0	0.338	1.587
A=121	N	1.28	45.5	0.413	1.613
Z=49	Z	1.24	55.1	0.341	1.587
A=127	N	1.28	43.8	0.413	1.613
Z=53	Z	1.24	59.8	0.350	1.587

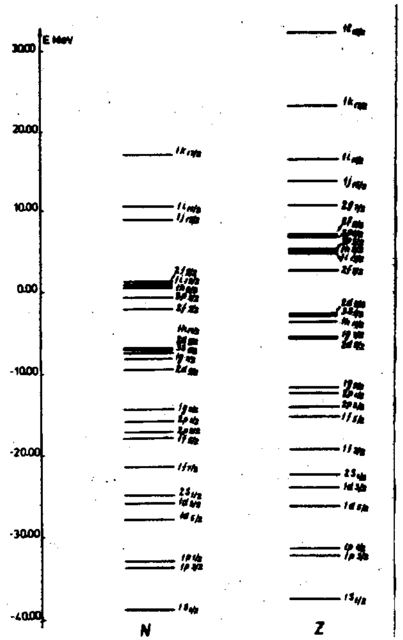

Fig. 1. Neutron and proton single-particle schemes for ^{124}Te (the Saxon-Woods potential parameters are given in table 1).

shows the single-particle neutron and proton levels for the scheme with $A=127$ used in the calculation for ^{124}Te. The proton scheme takes into account all the levels up to an energy of 15 MeV. The neutron scheme includes all the levels from the shell which is next to the unfilled one for the exception of the $3p_{1/2}$ subshell in ^{120}Sn and ^{124}Te. In the neutron scheme for ^{90}Zr the subshells $3p_{3/2}$ and $2f_{5/2}$ are also absent. At higher excitation energies (7-15 MeV) the subshells $2g_{9/2}$, $2g_{7/2}$, $3d_{5/2}$, $3d_{3/2}$ and $4s_{1/2}$ are absent. Thus, we lose mainly the levels with small spins because

of a small centrifugal barrier. From the figures of ref. [25] representing the matrix elements forming the giant resonances we conclude that in our calculations all the large matrix elements for the El resonance are taken into account and for the isoscalar E2 resonance only one large mattix element between neutron subshells $2g_{9/2}$ and $2d_{5/2}$ is lost. Thus, the matrix elements we have lost are small enough and should not influence significantly the results. In the isovector quadrupole resonance range several one-phonon states are lost. Since the single-particle levels far from the Fermi level take part in the formation of higher multipolarity resonances we calculate only the low-energy part (till 10 MeV) of the octupole resonance.

The superconducting pairing constants are fitted to experimental pairing energies [27]. The constants of the quadrupole isoscalar and isovector forces are fitted to the 2_1^+ level energy and the position of the isovector resonance. We get $k_2 = \kappa_1^2 / \kappa_0^2 \approx (-1.5-2.0)$. Varying of κ_1^2 little influences the 2_1^+ level and isoscalar E2 resonance characteristics. The constants κ_0^3 and κ_1^3 are taken to obtain the 3_1^- level energy, the ratio $k_3 = \kappa_1^3/\kappa_0^3 = -4.5$ being fixed according to the estimate of ref. [5]. For the deformed nuclei [8] $k_3 = -1.5$ which corresponds to the hydrodynamical model our calculations show that the RPA distribution B(E3) up to energies ~10 MeV coincides for both the k_3 values.

In describing the dipole resonance we should exclude the spurious state due to the center-mass motion of the nucleus. To this end, we employ the method of ref. [30]. According to it the Hamiltonian is supplemented with isoscalar dipole forces with a constant κ_0^1, which is chosen such that the first root of eq. (2) for $\lambda=1$ is zero. As is shown in ref. [31] this method, gives a 90% concentration of the spurious state on the root with $\omega=0$. If the spurious state is not excluded, the spurious admixtures change the structure of 8-12 MeV 1^- states and weakly affect the one-phonon states near the giant dipole resonance. The isovector dipole constant κ_1^1 is fitted to the experimental position of the E1 resonance. In our calculation we take into account the $\lambda \leq 8$ phonons. The constants for them are chosen so small that the one-phonon state structure is close to the corresponding quasiparticle state structure.

In our approach we treat phonons as boson neglecting their fermion structure and to a certain extent violate the Pauli principle when constructing the two-phonon components. A strict elimination of the relevant effects is very difficult to be realized due to a very large number of one- and two-phonon states used in our calculations. Nevertheless, to weaken possible distortions we introduce the following limitations to the wave function (5). One of the phonons in its two-phonon part is to be collective. As a result, the function (5) completely loses the components consisting of two noncollective phonons which may just violate the Pauli principle. Besides, since the collective phonons possess either very low (~ 1-2 MeV) or very high (~ 12-16 MeV) excitation energy, the

two-phonon states in the range of the isoscalar E2 or isovector E1 resonance are formed by the phonons of different quasiparticle structure. This also contributes to weakening of the Pauli principle violation.

4. THE GIANT DIPOLE RESONANCE

We start the discussion of the results for the giant dipole resonance with the nucleus ^{90}Zr. This nucleus has most thoroughly been studied by both experimenters [32,33] and theorists [7,11]. In RPA the E1 resonance of this nucleus exhibits, according to our calculations, 7 collective states in the energy range 14.8-18.0 MeV (fig. 2a). The figure shows the states which give a contribution equal to or more than 1 per cent to the EWSR, the most collective of them being of a 16 MeV energy. Table 2 gives the total contribution of one-phonon states in specified energy ranges to the EWSR. It is easily seen that in the range 14.0-18.0 MeV the strength exhausts 71% of the EWSR. The state with maximum B(E1) values gives a 38.6% contribution. Up to 30 MeV, about 90% of the EWSR are found to be exhausted which evidences for a rather large space of a single-particle states. These results are in satisfactory agreement with the calculations of Liu and Brown [11].

The results of calculations of the E1 resonance taking into account the quasiparticle-phonon interaction are shown in fig. 2b. The calculation has been performed with $\Delta = 0.001$ MeV, and, therefore, the

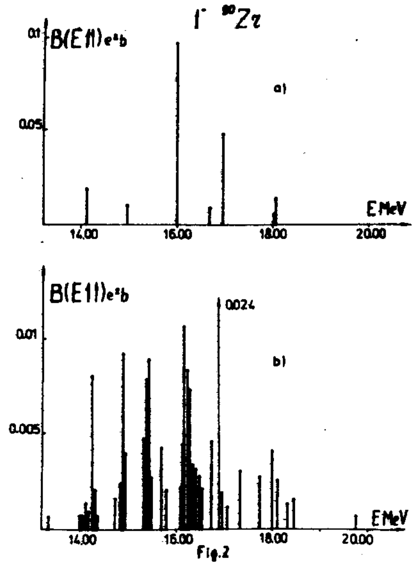

Fig. 2. a. Dipole isovector resonance in ^{90}Zr in RPA. b. Dipole isovector resonance in ^{90}Zr, calculated with the account of the qiasiparticle-phonon interaction ($\Delta = 0.001$ MeV).

Table 2
Contribution of one-phonon 1^- states to the EWSR

Energy interval, MeV	^{90}Zr % EWSR	^{120}Sn % EWSR	^{124}Te %EWSR
0-10	1.18	1.08	1.10
10-14	4.55	9.74	10.81
14-16	47.25	65.8	4.89
16-18	24.0	11.3	70.4
18-20	9.04	0.0	4.45
20-25	1.57	1.43	3.80
25-30	1.40	3.20	2.89
0-30	89.5	92.5	98.3

results are presented as separate solutions. The amplitude of each peak of fig. 2a is obtained as $\int_{\Delta E} b(E1,\eta)\,d\eta$ where ΔE is the energy interval in which the peak is localized (as far as Δ is very small, the peaks do not overlap). In the one-phonon part of the wave function (5) only the collective states of fig. 2a are summed up, that is, the rank of the determinant (6) is 7.

Thus, the one-phonon component strength has been distributed over near lying two-phonon states. As a result, the transition from individual states has become less intensive, but the number of states with noticeable $B(E1, 0^+_{g.s.} \to 1^-_\nu)$ values has become much larger. The giant dipole resonance forms now several dozens of states in the interval 13.5-18.5 MeV. In this excitation region the strength exhausts 63% of the EWSR. In RPA the states lying in this interval give a 72% contribution to the EWSR. Thus, the quasiparticle-phonon interaction makes the concentration of the sipole transition strength less strong. It is interesting to note that the maximum in the $B(E1)$ distribution in fig. 2b is shifted up by about 1 MeV with respect to the RPA meximum. This is due to different coupling strength of the two most collective states in fig. 2a with two-phonon configurations. As a result, the state with larger $B(E1)$ value has been distributed more strongly over the roots eq. (6).

We note that the wave function (5) contains two-phonon components which can be expressed in terms of the 2p–2h configurations. Therefore, it contains some admixture of the components with isotopic spin $T_> = T_< + 1$. On the basis of our calculations one may try to find the isotopic splitting of the isovector dipole resonance to which it was indicated in refs./33,34/.

We go over to the discussion of the results for ^{120}Sn (fig. 3). In RPA the giant dipole resonance in ^{120}Sn is obtained as

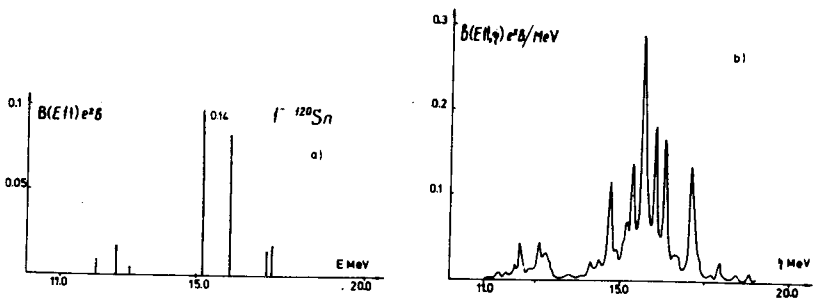

Fig. 3. a. Dipole isovector resonance in ^{120}Sn in RPA. b. Strength function $b(E1,\eta)$ in ^{120}Sn, calculated with the account of the quasiparticle-phonon interaction ($\Delta = 0.1$ MeV)

three groups of collective states lying in the intervals 12-13 MeV, 15-16 MeV and 17-17.5 MeV. Their total cobtribution to the EWSR is 83%, the two peaks at 15.5 and 16.0 MeV exhaust 66%. As is seen from table 2, up to 30 MeV it is exhausted more

than 90% of the EWSR which gives an additional evidence for the completeness of our basis.

The function $b(E1,\eta)$ is calculated with $\Delta = 0.1$ MeV, the rank of the determinant (6) is 7. The maximum distribution is got at 15.7 MeV, i.e. between the two most collective solutions in RPA. This is also in agreement with the experimental data of refs. /32,35/. In the region of 12-18 MeV excitation the strength exhausts 73% of the EWSR, however, the largest portion of the dipole transition strength (68%) is concentrated in the interval 14-18 MeV. This value is slightly less than the experimental resonance width $\Gamma = 5$ MeV /32,35/.

The main difference of the ^{124}Te nucleus from the ^{90}Zr and ^{120}Sn nuclei consists in that the former is not magic by the proton nor by the neutron number. As low-lying state calculations show in ^{124}Te the quasiparticle-phonon interaction is much stronger. This is also seen from the results obtained by us for giant resonances in this nucleus.

The RPA calculations (fig. 4a) show that the giant E1 resonance in ^{124}Te is formed by a large number of collective states in the region of 12-19 MeV. The one-phonon 1^- states of an energy up to 30 MeV in ^{124}Te exhaust 98%, the collective states in the interval 12-20 MeV - 84%, and the most collective of them - 57% of the EWSR.

The quasiparticle-phonon interaction results in a noticeable broadening of the resonance, as is seen from fig. 4b. The calculation is performed with $\Delta = 0.4$ MeV, and in the one-phonon part of the function (5)

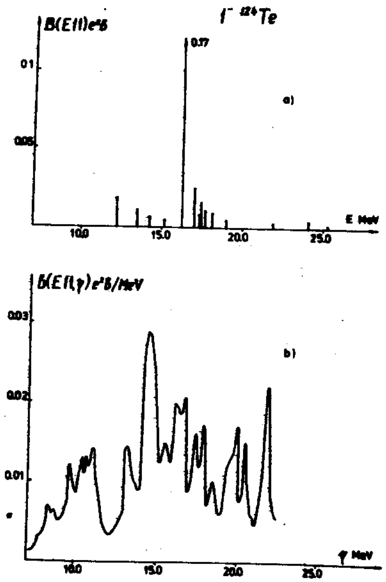

Fig. 4. a. Dipole isovector resonance in ^{124}Te in RPA. b. Strength function $b(E1,\eta)$ in ^{124}Te calculated with the account of the quasi-particle-phonon interaction ($\Delta = 0.4$ MeV).

ten most collective roots are taken into account (i.e., the phonon of an energy of 19 MeV which gives a 1.5% contribution to the EWSR is disregarded). The maximum in the B(E1) distribution drops down to an energy of about 14.5 MeV which is somewhat lower than the experimental value /37/. Now between 12.0 and 20 MeV it is exhausted only 31.5% of the EWSR, i.e., a considerable portion of the dipole transition strength was displaced toward other excitation energy regions. So, the 20-22 MeV interval gives a 8.1% contribution to the

EWSR while in RPA this contribution is 1.5%. Thus, the situation in ^{124}Te differs strongly from that in both ^{90}Zr and ^{120}Sn nuclei where the quasiparticle-phonon interaction redistribute the dipole electric transition strength along the energy scale not so strongly. We note that in experiments too, the total photoneutron cross section at 20-24 MeV in ^{124}Te ref. /37/ is seen to be higher than that in ^{120}Sn (ref. /35/).

5. THE GIANT QUADRUPOLE RESONANCE

The giant isoscalar quadrupole resonance has experimentally been studied much better than other "new" resonances. The experimental information about the isovector quadrupole resonance is still very poor /1-3/. So, of three nuclei, we are interested in in this paper, some experimental evidence for the existence of an isovector quadrupole resonance is available only for ^{90}Zr refs. /38-41/.

The results of calculations in RPA for ^{90}Zr are given in fig.5.a. As before, the states giving a 1% contribution to the EWSR are presented. The isoscalar resonance is formed by two- one-phonon 14 MeV states spaced by 0.3 MeV. The contribution of these states to the EWSR is 26%. The isovector quadrupole resonance is formed by several states lying in the range of 25-28 MeV, their contribution to the EWSR is 39%. A number of the one-phonon states is seen to be lost due to incompleteness of the single-particle basis. This is also proved by the degree to which the EWSR strength

Table 3

Contribution of one-phonon 2^+ states to the EWSR

Energy interval, MeV	^{90}Zr % EWSR	^{120}Sn % EWSR	^{124}Te % EWSR
0÷10	5.05	8.61	10.1
10÷12	0.25	1.25	3.16
12÷14	26.2	26.2	24.9
14÷16	0.54	0.19	0.15
16÷18	0.57	3.41	3.91
18÷20	3.30	2.57	5.36
20÷25	3.48	17.6	35.3
25÷30	41.72	18.4	9.43
0÷30	81.1	78.6	92.3

is exhausted (see table 3). The one-phonon 2^+ states up to an excitation energy of 30 MeV exhaust 81% of the EWSR. All these data are in a rather satisfactory agreement with the experimental values [38,39,41] and the theoretical calculations of Liu and Brown [11]. However, the contribution of the isovector resonance to the EWSR obtained by us is essentially greater than that of the authors of ref. [11].

The calculation which accountes for the quasiparticle-phonon interaction is performed in the same manner as for the dipole resonance with small Δ and, therefore, the results of fig. 5b are presented as individual states the contribution of which to the EWSR is greater than 0.2%. Now the isoscalar resonance is formed by a large group of states concentrated in the region

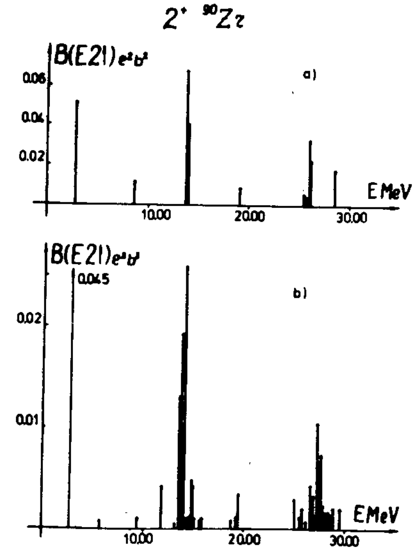

Fig. 5. a. Collective 2^+ states in ^{90}Zr, calculated in RPA in the excitation energy interval 0-30 MeV; b. Collective 2^+ states in ^{90}Zr calculated with the account of the quasiparticle-phonon interaction (Δ = = 0.001 MeV).

of 12-16 MeV. The experimental resonance width [38,41] is somewhat larger than 4 MeV. A similar picture holds for the isovector resonance. It occupies now the region 25-30 MeV and gives to the EWSR a 31% contribution. When calculating in the isoscalar resonance region we have taken into account 7 most collective quadrupole states from the region between zero and 30 MeV. When calculating in the isovector resonance region collective states of an energy lower than 10.0 MeV have not been included in the one-phonon part of the wave function (5), but all the collective states forming the isovector resonance in RPA have been taken into account.

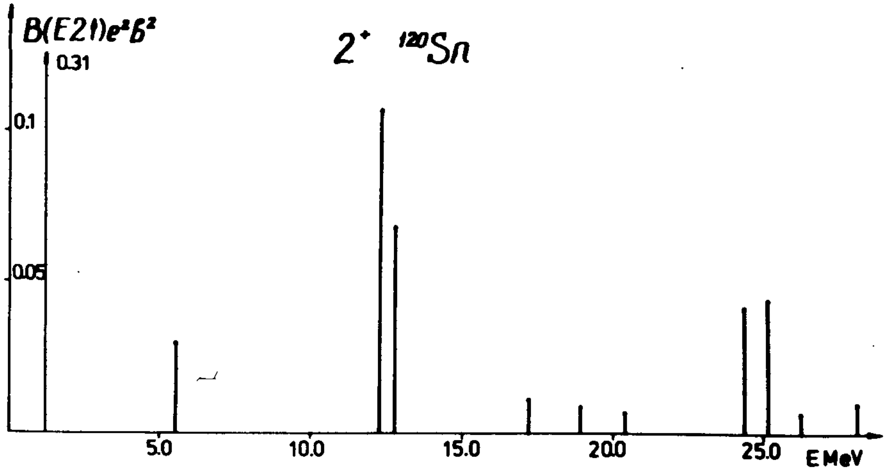

Fig. 6. Collective 2^+ states in ^{120}Sn in the excitation energy interval 0-30 MeV, calculated in RPA.

Isovector quadrupole resonances have not been observed in ^{120}Sn nor in ^{124}Te. Therefore in these nuclei we have restricted ourselves to the calculations accounting for the quasiparticle-phonon interaction only for isoscalar modes, however, the isovector quadrupole forces have been included in the Hamiltonian. The picture obtained for the collective 2^+ excitations in ^{120}Sn is given in fig. 6 . In this case too, the isoscalar resonance is formed by only two one-phonon states of energies 12.2 and 12.8 MeV, respectively. They give to the EWSR a 26.5% contribution, that is the same as in ^{90}Zr. In the nucleus ^{120}Sn, up to an energy of 30 MeV, it is exhausted about 79% of the EWSR, i.e., in this case too, we lose a number of one-phonon states. However, as was discussed in Section 3, the

main losses occur in the isovector resonance region. Due to the quasiparticle-phonon interaction the isoscalar quadrupole resonance becomes broader (fig. 8), the maximum in the B(E2) distribution is slighly displaced in favor of lower energies. As a result, in the interval 12.0-14.0 MeV it is now exhausted only 11.5% of the EWSR, and 25% contribution to the EWSR comes from the interval 10-15 MeV.

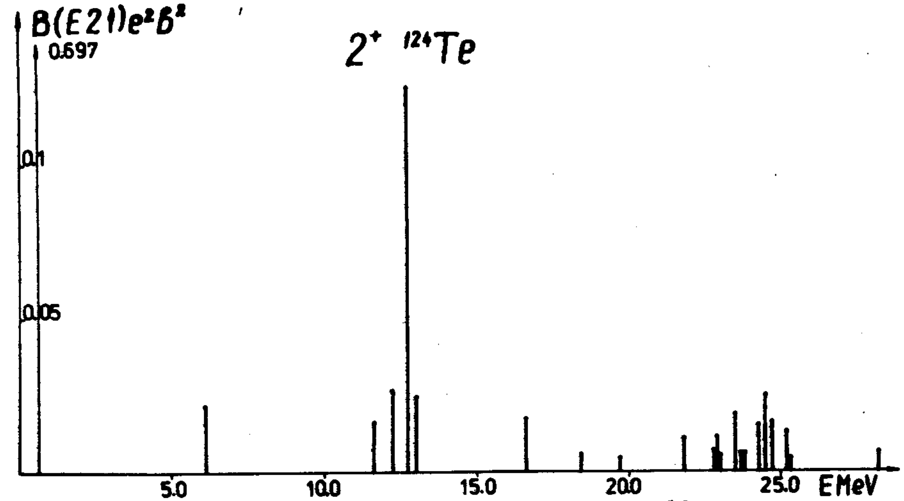

Fig. 7. Collective 2^+ states in ^{124}Te at 0-30 MeV energy calculated in RPA.

Even stronger changes are due to the quasiparticle-phonon interaction in ^{124}Te. The four collective states forming the isoscalar resonance in RPA (fig. 7) lie at energies 11.0-14.0 MeV and give a contribution of 26%, all the states exhaust 93% of the EWSR (table 3). After the interaction with two-phonon configurations has been included, the region in which there exist

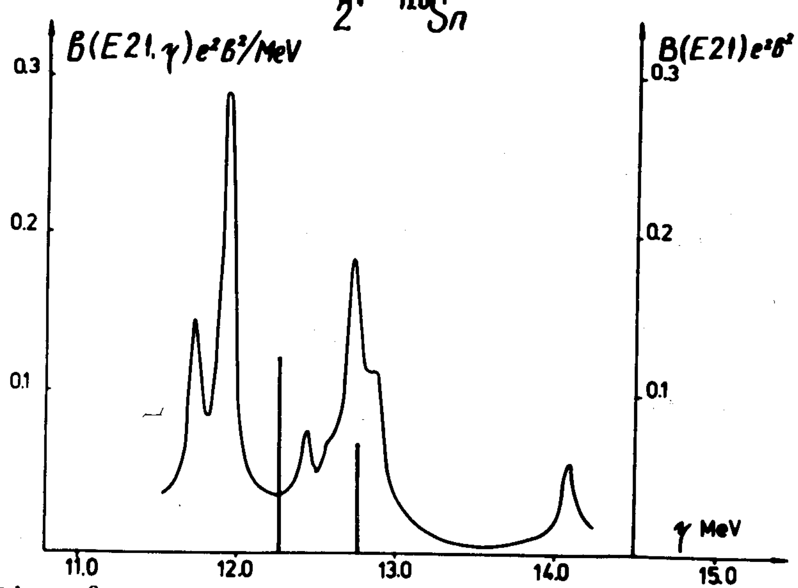

Fig. 8. Strength function $b(E2, \eta)$ in ^{120}Sn calculated with the account of the quasi-particle-phonon interaction in the region of isoscalar quadrupole resonance ($\Delta = 0.1$ MeV). The B(E2) values for collective quadrupole RPA states are shown by the vertical lines (see also fig. 6).

noticeable E2 transitions is seen to increase up to 10 MeV (6-16 MeV), fig. 9 (as in the case of the E1 resonance the calculation has been performed with $\Delta = 0.4$ MeV). In the 6-16 MeV interval it is exhausted 25.6% of the EWSR, the region of the $b(E2,\eta)$ maximum, 13.5-16.5 MeV, gives to the EWSR a 9% contribution. Thus by the example of the E1 and E2 resonances in ^{124}Te it is seen that the quasiparticle-phonon interaction in those nuclei, where the low-lying vibrational states are strongly collectivized, greatly influences high excitations, too.

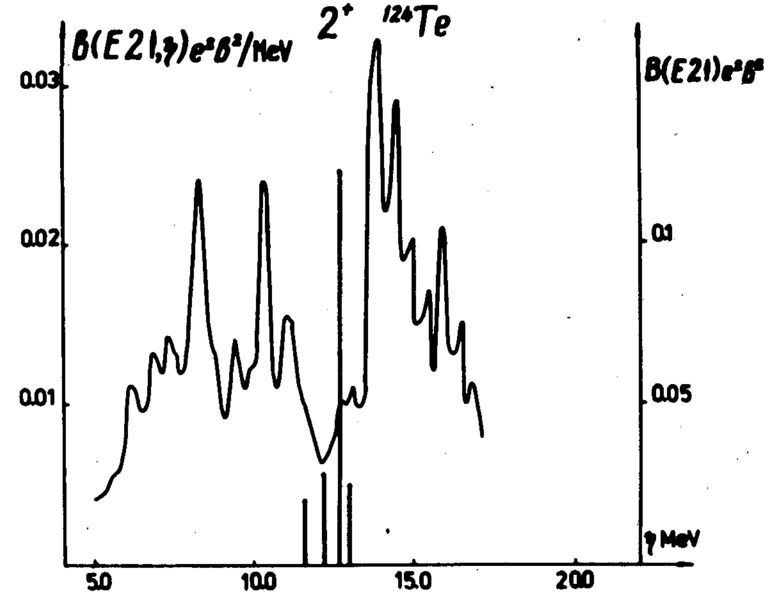

Fig. 9. Strength function $b(E2, \eta)$ in ^{124}Te calculated taking into account the quasi-particle-phonon interaction in the region of isoscalar quadrupole resonance ($\Delta = 0.4$ MeV). The notation is the same as in fig. 8.

6. LOW-ENERGY OCTUPOLE RESONANCES

The giant-resonance-like structures have been observed in ref. /5/ at $E_{exc} \sim 32 A^{-1/3}$ in nuclei from ^{90}Zr to ^{154}Sm. Analysis of the angular distributions had led to an assignment of $J^{\pi} = 3^-$ and an isoscalar EWSR (eq. (15)) fraction of 16-22% for this structure. The existence of similar structures was also indicated in theoretical studies for both the spherical /11/ and deformed /8/ nuclei.

Most detailed information about the octupole excitations at energies $\lesssim 9$ MeV has been extracted for ^{90}Zr.[5] Up to an energy of 8 MeV it was observed 6 levels with $J^{\pi}=3^-$ and a $\sim 2 \div 3$ s.p.u.B(E3) value. Up to 9 MeV in ^{90}Zr it is exhausted 26% of the EWSR, the 3_1^- state exhausting 7%.

We have calculated the 3^- states in ^{90}Zr up to an energy of 10 MeV both in RPA and taking into account two-phonon admixtures. In RPA, up to 8 MeV, in ^{90}Zr we have 6 one-phonon states, three of them (very collective) are shown in fig. 10 by the dashed lines. The excitation probabilities for the three remaining ones are by a factor of $10^2 - 10^3$ smaller than those presented in the figure. Up to 9 MeV it is

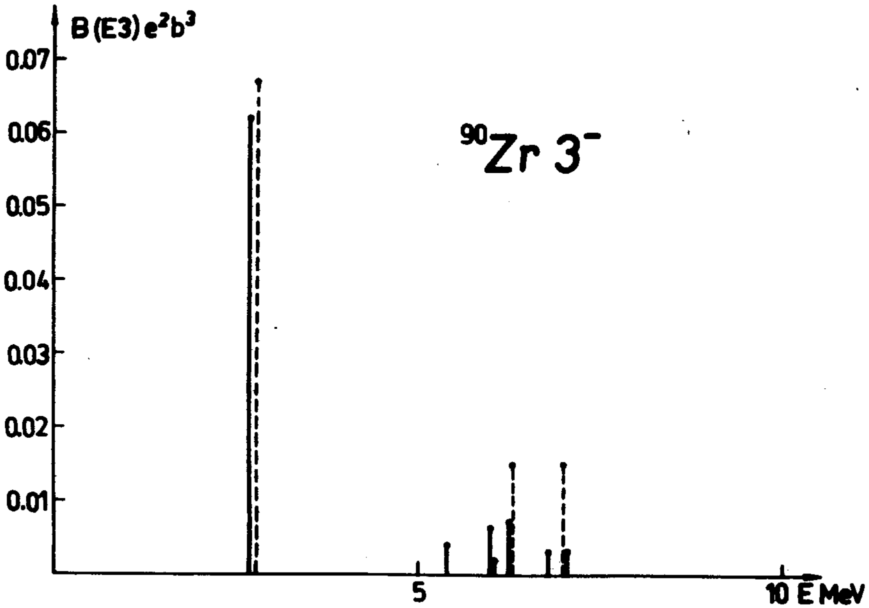

Fig. 10. Collective 3^- states in ^{90}Zr in the energy interval 0-30 MeV. Dotted lines is the calculation in RPA, continuous lines is the calculation taking into account the quasiparticle-phonon interaction.

exhausted 25.8% of the isoscalar EWSR (for three collective roots - 25.2% and the 3_1^- states - 12%). Thus, the low-energy octupole resonance (LEOR) in RPA together with the octupole-octupole residual interaction form two collective 3^- states at 6.5-7.0 MeV giving a 13.2% contribution to the isoscalar EWSR.

After including the quasiparticle-phonon interaction we have a picture presented in fig. 10 by continuous lines. The calculation is performed as for ^{90}Zr with $\Delta = 0.001$ MeV which enables individual solutions to be separated. The one-phonon part of the wave function (5) includes only three above-mentioned collective states. The interaction between one- and two-phonon states has resulted in a further distribution of the E3 transition strength. The 3_1^- state has little been affected by the quasiparticle-phonon interaction: its energy has been lowered by only 0.1 MeV and the B(E3) value has slightly been decreased. As a result, the 3_1^- state contribution to the isoscalar EWSR has decreased down to 10.8%. The states forming LEOR have been splitted and now it is concentrated in the range 5.4-7.1 MeV and consists of six states with B(E3)~1 s.p.u. The energy position of these states and their excitation probabilities are in agreement, in the main features, with experiment[5]. The contribution of the octupole states up to an energy of 9 MeV to the isoscalar EWSR remains actually the same and amounts to 24.9%.

For ^{120}Sn there is no experimental information about LEOR. Ref.[5] gives some data on LEOR for ^{118}Sn. The LEOR energy in this

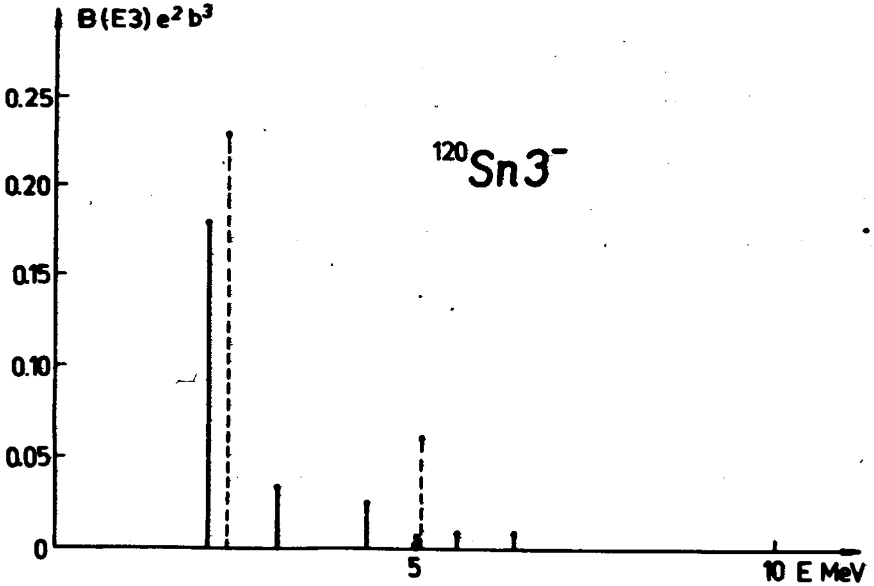

Fig. 11. Collective 3^- states in ^{120}Sn in the energy interval 0-30 MeV. The notation is the same as in Fig. 10.

nucleus is about 6.9 MeV and their contribution to the isoscalar EWSR is about 20%. In ^{118}Sn in the interval 0-8 MeV it is exhausted about 30% of the isoscalar EWSR. This picture is similar to our calculations in ^{120}Sn (fig. 11). In our case, in RPA (dashed lines) the states up to 8 MeV exhaust 35% of the isoscalar EWSR. Two of these states are collective: the 3^-_1 level (20.5%) and the E=5.1 level (about 12%). The second level is just expected to correspond to LEOR, but is located by about 2 MeV lower. The interaction with the two-phonon configurations distributes B(E3) over many

levels (although for each level B(E3) is naturally smaller). The collective 3^- levels are almost uniformly distributed in the interval 2.0-6.5 MeV and do not form, as in ^{90}Zr, the group of levels removed from the 3^-_1 state and strictly localized. The states presented in fig. 11 by continuous lines exhaust 32.2% for the isoscalar EWSR.

7. CONCLUSION

We have calculated the B(Eλ) values in RPA and the strength functions $b(E\lambda,\eta)$ taking into account the quasiparticle-phonon interaction for giant multipole resonances in spherical double even nuclei. By the example of the ^{90}Zr nucleus, it is seen that our results are in satisfactory agreement with those observed and predicted by other authors. It is interesting to note that although in ref. /11/ the calculation was performed with a large configuration space in RPA based on the Hartree-Fock ground states of newly developed Skyrme interaction and in our we used a simple multipole-multipole interaction, the positions and strengths of the giant multipole resonance transitions in ^{90}Zr were found to be close to each other in both calculations.

Long ago /13-15/ it was indicated that the distribution of the dipole strength in a wide energy interval is connected with the interaction of the giant dipole resonance with low-lying collective states of nuclei. However, so far the account of this interaction in medium and heavy nuclei has been made only phenomenologically. In the present paper it is shown that the model accounting for the quasiparticle-phonon interaction is valid for the description of the photoproduction strength function for the giant multipole resonance. The quasi-particle-phonon interaction results in a strong fragmentation of the one-phonon states forming giant multipole resonances in RPA. The electromagnetic transition strength distribution over two-phonon states enables one to obtain roughly the position and the magnitude of the energy regions where the giant multipole resonances are localized.

The giant resonance characteristics are most strongly affected by the quasipar-ticle-phonon interaction in nuclei with strongly collectivized low-energy excita-tions, such as ^{124}Te. In nuclei, like ^{90}Zr and ^{120}Sn, this interaction is not very im-portant. A further complication of the wave function (5) by including three-phonon com-ponents may well lead to some smoothing of the function $b(E\lambda,\eta)$. For example, this may result in a disappearance of the dip at 12 MeV for the E1 resonance in ^{124}Te. How-ever, it seems to us that the intro-duction of three-phonon components does not lead to a significant change of the results obtained.

We would like to stress that our model does not pretend to a correct description of the wave functions of highly excitated states. In ref. [42] it is demonstrated that our model may serve as a basis for the description of the strength distribution of only few quasiparticle components at low, intermediate and high excitation energies. We describe the excitation probabilities for a set of states rather than for individual states in fixed energy intervals. As far as the total photoexcitation cross sections from the ground states of double even nuclei are mainly defined by the one-phonon fragmentation, our model is found to be suitable for the description of the strength functions $b(E\lambda,\eta)$ for giant multipole reso-nances. These calculations are a new step towards studying giant multipole resonances compared with the RPA calculations.

It should be noted that the main contri-bution to the wave functions of highly excited states comes from many-quasiparticle com-ponents. It is undoutful that in future new properties of highly excited states due to many-quasiparticle components will be demonstrated. So far, there is no data on the magnitude and distribution of many-quasiparticle components of the highly excited state wave functions. Even for neut-ron resonances it is shown [43] that direct experimental evidence for many-quasipar-ticle components of their wave functions is absent. The contribution of the few-quasi-particle components to the wave function normalization is only 10^{-4} - 10^{-6}.

Thus, the model based on the quasipar-ticle-phonon interaction gives the general description for the strength functions in one-nucleon transfer reactions at inter-mediate excitation energies, the neutron strength functions-at neutron binding ener-gy and for the strength functions $b(E\lambda,\eta)$- in the region of giant multipole resonances.

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