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DESCRIPTION OF $T_>$ GIANT RESONANCES IN SPHERICAL NUCLEI

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Introduction

The few-quasiparticle components of the nuclear state wave functions have been calculated within the quasiparticle-phonon nuclear model [1] at low, intermediate and high excitation energies. The energies and wave functions of the low-lying nonrotational states in deformed nuclei have been calculated within this model [2,3]. The fragmentation of one- and two-quasiparticle states has been calculated and the neutron s- and p-wave strength functions in spherical [4-7] and deformed [8] nuclei have been obtained. The energies and widths of giant multipole and spin-multipole resonances in spherical [4,9-12] and deformed [12-14] nuclei have been described.

In recent years much attention has been paid to the study of the $T_>$ giant resonances ($T_> = T_0 + 1$, T_0 is the isospin of the ground state), Gamov-Teller 1^+ resonances, first forbidden 1^- charge exchange resonances, analog hole states and other characteristics of excited states [15,16]. More experimental information is accumulated on the $T_>$ giant dipole resonance in spherical nuclei. Several theoretical approaches exist for the description of analog states and charge-exchange giant resonances [17-21].

Within the quasiparticle-phonon nuclear model one can describe the isobar-analog states, $T_>$ giant isovector resonances, charge-exchange resonances and others. In this paper we present the mathematical apparatus for describing the $T_>$ giant resonances and the results of calculations for the energies and B(E1)-values of the $T_>$ giant dipole resonances in several spherical nuclei.

1. The model

The Hamiltonian of the quasiparticle-phonon nuclear model includes an average field as the Saxon-Woods potential and the superconducting paiting correlations. It also contains the multipole-multipole and spin-multipole - spin-multipole isoscalar and isovector and isovector interactions in the particle-hole and particle-particle channels. The multipole-multipole interaction is

$$V_\lambda(\vec{r}_1,\vec{r}_2) = -\left(\frac{1}{2}\varkappa_0^{(\lambda)} + 2\varkappa_1^{(\lambda)}\vec{t}_1\vec{t}_2\right) r_1^\lambda Y_{\lambda\mu}(\theta_1,\varphi_1)\, r_2^\lambda Y_{\lambda,-\mu}(\theta_2,\varphi_2). \qquad (1)$$

To describe the $T_>$ giant multipole resonances, one should introduce n-p (neutron-proton) phonons generated by the following part of interaction (1):

$$\varkappa_1^{(\lambda)}\left(t_1^{(+)}t_2^{(-)} + t_1^{(-)}t_2^{(+)}\right) r_1^\lambda Y_{\lambda\mu}(\theta_1,\varphi_1)\, r_2^\lambda Y_{\lambda,-\mu}(\theta_2,\varphi_2). \qquad (2)$$

The corresponding part of the Hamiltonian will be

$$H_{coll}^{np} = -2\sum_{\lambda,\mu}\varkappa_1^{(\lambda)}\Omega^+_{\lambda\mu}\Omega_{\lambda\mu}. \qquad (3)$$

$$\Omega_{\lambda\mu} = \sum_{\substack{j_p,m_p\\ j_n,m_n}} \frac{(-1)^{j_n-m_n}}{\sqrt{2\lambda+1}} \langle j_p m_p j_n -m_n|\lambda\mu\rangle\langle j_p\| i^\lambda r^\lambda Y_\lambda t^{(-)}\|j_n\rangle \cdot a^+_{j_p m_p}\, a_{j_n m_n} \qquad (4)$$

Here the single-particle states are specified by the quantum numbers: proton $j_p m_p$, neutron $j_n m_n$; a^+_{jm} and a_{jm} are the nucleon creation and absorption operators.

Now we perform the Bogolubov canonical transformation, introduce the operators

$$A(j_p j_n;\lambda\mu) = \sum_{m_p m_n}\langle j_p m_p j_n m_n|\lambda\mu\rangle \alpha_{j_n m_n}\alpha_{j_p m_p},$$

$$B(j_p j_n;\lambda\mu) = \sum_{m_p m_n}(-1)^{j_n+m_n}\langle j_p m_p j_n m_n|\lambda\mu\rangle \alpha^+_{j_p m_p}\alpha_{j_n,-m_n},$$

where α^+_{jm} and α_{jm} are the quasiparticle creation and absorption operators, and get

$$\Omega_{\lambda\mu} = \sum_{j_p m_p}\frac{f^\lambda(j_p j_n)}{\sqrt{2\lambda+1}}\Big\{u_{j_p}v_{j_n}A^+(j_p j_n;\lambda\mu) + (-1)^{\lambda-\mu}v_{j_p}u_{j_n}A(j_p j_n;\lambda-\mu) + u_{j_p}u_{j_n}B(j_p j_n;\lambda\mu) + (-1)^{j_p+j_n-\lambda}v_{j_p}v_{j_n}B(j_n j_p;\lambda-\mu)\Big\}, \qquad (4')$$

where $f^\lambda(j_p j_n) = \langle j_p\| i^\lambda r^\lambda Y_\lambda t^{(-)}\|j_n\rangle$, u_j, v and v_j are the Bogolubov transformation coefficients. Now we introduce the operators of n-p phonons

$$\Omega_{\lambda\mu i} = \frac{1}{\sqrt{2}}\sum_{j_p j_n}\Big\{\psi^{\lambda i}_{j_p j_n}A(j_p j_n;\lambda\mu) - (-1)^{\lambda-\mu}\varphi^{\lambda i}_{j_p j_n}A^+(j_p j_n;\lambda-\mu)\Big\}. \qquad (5)$$

In the quasiboson approximation they satisfy the condition

$$[\Omega_{\lambda'\mu' i'},\Omega^+_{\lambda\mu i}] = \frac{1}{2}\sum_{j_p j_n}\Big\{\psi^{\lambda' i'}_{j_p j_n}\psi^{\lambda i}_{j_p j_n} - \varphi^{\lambda' i'}_{j_p j_n}\varphi^{\lambda i}_{j_p j_n}\Big\} = \delta_{\lambda\lambda'}\delta_{\mu\mu'}\delta_{ii'}.$$

To describe n-p phonons, we use the following part of the model Hamiltonian (6)

$$H_v^{np} = \sum_{j_p}\varepsilon(j_p)\sum_{m_p}\alpha^+_{j_p m_p}\alpha_{j_p m_p} + \sum_{j_n}\varepsilon(j_n)\sum_{m_n}\alpha^+_{j_n m_n}\alpha_{j_n m_n} -$$

$$-\sum_{\lambda,\mu,i,i'} \frac{\varkappa_1^{(\lambda)}}{2\lambda+1} \sum_{j_p j_n j_p' j_n'} f^{\lambda}(j_p,j_n)\, f^{\lambda}(j_n',j_p') \cdot \tag{7}$$

$$\cdot\Big\{\big(\psi^{\lambda i}_{j_pj_n} v_{j_p} u_{j_n} + \varphi^{\lambda i}_{j_pj_n} u_{j_p} v_{j_n}\big)\big(\psi^{\lambda i'}_{j_p'j_n'} v_{j_p'} u_{j_n'} + \varphi^{\lambda i'}_{j_p'j_n'} u_{j_p'} v_{j_n'}\big) +$$

$$+\big(\psi^{\lambda i}_{j_pj_n} u_{j_p} v_{j_n} + \varphi^{\lambda i}_{j_pj_n} v_{j_p} u_{j_n}\big)\big(\psi^{\lambda i'}_{j_p'j_n'} u_{j_p'} v_{j_n'} + \varphi^{\lambda i'}_{j_p'j_n'} v_{j_p'} u_{j_n'}\big)\Big\}\Omega^+_{\lambda\mu i}\Omega_{\lambda\mu i'}.$$

We introduce the wave function of the one-phonon state of an odd-odd nucleus

$$\Omega^+_{\lambda\mu i}|\rangle, \tag{8}$$

where $|\rangle$ is the ground state wave function of a doubly even nucleus satisfying the condition $\Omega_{\lambda\mu i}|\rangle = 0$.

The energies of the one-phonon states $\omega_{\lambda i}$ are obtained by the variational principle

$$\delta\Big\{\langle|\Omega_{\lambda\mu i} H_v^{np}\Omega^+_{\lambda\mu i}|\rangle - \frac{\omega_{\lambda i}}{2}\Big[\sum_{j_pj_n}\{(\psi^{\lambda i}_{j_pj_n})^2-(\varphi^{\lambda i}_{j_pj_n})^2\}-2\Big]\Big\}=0 \tag{9}$$

After transformation the secular equation for finding the energies $\omega_{\lambda i}$ is

$$\mathcal{F}(\omega_{\lambda i}) = (1-\varkappa_1^{(\lambda)}X_{11}^{\lambda i})(1-\varkappa_1^{(\lambda)}X_{22}^{\lambda i}) - (\varkappa_1^{(\lambda)})^2 (X_{12}^{\lambda i})^2 = 0, \tag{10}$$

where

$$X_{11}^{\lambda i} = \frac{2}{2\lambda+1}\sum_{j_pj_n}|f^\lambda(j_pj_n)|^2\Big\{\frac{u_{j_p}^2 v_{j_n}^2}{\varepsilon(j_pj_n)-\omega_{\lambda i}} + \frac{v_{j_p}^2 u_{j_n}^2}{\varepsilon(j_pj_n)+\omega_{\lambda i}}\Big\},$$

$$X_{12}^{\lambda i} = \frac{2}{2\lambda+1}\sum_{j_pj_n}|f^\lambda(j_pj_n)|^2\Big\{\frac{u_{j_p}u_{j_n}v_{j_p}v_{j_n}}{\varepsilon(j_pj_n)-\omega_{\lambda i}} + \frac{u_{j_p}u_{j_n}v_{j_p}v_{j_n}}{\varepsilon(j_pj_n)+\omega_{\lambda i}}\Big\},$$

$$X_{22}^{\lambda i} = \frac{2}{2\lambda+1}\sum_{j_pj_n}|f^\lambda(j_pj_n)|^2\Big\{\frac{v_{j_p}^2 u_{j_n}^2}{\varepsilon(j_pj_n)-\omega_{\lambda i}} + \frac{u_{j_p}^2 v_{j_n}^2}{\varepsilon(j_pj_n)+\omega_{\lambda i}}\Big\}.$$

$\varepsilon(j_pj_n) = \varepsilon(j_p)+\varepsilon(j_n)$. If the pairing is absent either in the neutron or proton system, then $X_{12}^{\lambda i} = 0$ and eq. (10) is divided into two equations. Solving eq. (10) at fixed $\varkappa_1^{(\lambda)}$, we find the energies $\omega_{\lambda i}$ of the one-phonon states. Using condition (6), we get

$$\psi^{\lambda i}_{j_pj_n} = 2\sqrt{\frac{\varkappa_1^{(\lambda)}}{1-\varkappa_1^{(\lambda)}X_{11}^{\lambda i}}}\,\frac{1}{\sqrt{-\frac{\partial\mathcal F(\omega)}{\partial\omega}\Big|_{\omega=\omega_{\lambda i}}}}\,\frac{f^{\lambda}(j_pj_n)}{\sqrt{2\lambda+1}}\,\frac{u_{j_p}v_{j_n}\varkappa_1^{(\lambda)}X_{12}^{\lambda i}+v_{j_p}u_{j_n}(1-\varkappa_1^{(\lambda)}X_{11}^{\lambda i})}{\varepsilon(j_pj_n)-\omega_{\lambda i}}, \tag{11}$$

$$\varphi^{\lambda i}_{j_pj_n} = 2\sqrt{\frac{\varkappa_1^{(\lambda)}}{1-\varkappa_1^{(\lambda)}X_{11}^{\lambda i}}}\,\frac{1}{\sqrt{-\frac{\partial\mathcal F(\omega)}{\partial\omega}\Big|_{\omega=\omega_{\lambda i}}}}\,\frac{f^{\lambda}(j_pj_n)}{\sqrt{2\lambda+1}}\,\frac{u_{j_p}v_{j_n}(1-\varkappa_1^{(\lambda)}X_{11}^{\lambda i})+v_{j_p}u_{j_n}\varkappa_1^{(\lambda)}X_{12}^{\lambda i}}{\varepsilon(j_pj_n)+\omega_{\lambda i}}. \tag{11'}$$

The secular equation (10) has the same form as the equations describing the Gamov-Teller 1^+ excitations in odd-odd nuclei /19,22,23/. The same procedure is used for finding the secular equations for the spin-multipole n-p phonons.

2. Wave functions of $T_>$ states of giant resonances and $E\lambda$-transition probabilities

Let the wave function $|\rangle$ of the ground state of a doubly even nucleus have isospin T_0 and its projection T_0. Then $\Omega^+_{\lambda\mu i}|\rangle$ describes odd-odd nuclei with $T_0 \pm 1$. By shifting the neutron and proton chemical potentials λ_n λ_p, may be achieved the wave function $\Omega^+_{\lambda\mu i}|\rangle$ to describe the state with T_0+1, T_0+1. To pass to the states T_0+1, T_0 of the considered nucleus, we use the operator $T^{(-)} = \sum_{K=1}^{A} t_K^{(-)}$ (see ref. /18/). As a result, the wave function of the $T_>$ one-phonon states of multipolarity λ in a doubly even nucleus with the isospin projection equal to T_0 can be written as

$$\frac{1}{\sqrt{2T_0+2}} T^{(-)} \Omega^{+}_{\lambda\mu i} | \rangle . \tag{12}$$

The energies of the one-phonon $T_>$ states with $I=\lambda$ and $\pi=(-1)^{\lambda}$ are

$$\frac{1}{2T_0+2} \langle | \Omega_{\lambda\mu i} T^{(+)} H_{v}^{np} T^{(-)} \Omega^{+}_{\lambda\mu i} | \rangle = \omega_{\lambda i} + \Delta E_c \tag{13}$$

since according to the formal theory of analog states/17/, the Coulomb energy ΔE_c is

$$\Delta E_c = \frac{\langle | [T^{(+)}, [H, T^{(-)}]] | \rangle}{\langle | T^{(+)} T^{(-)} | \rangle} . \tag{14}$$

The reduced $E\lambda$-transition probability from the ground state of a doubly nucleus to the state described by the wave function (12) is

$$B(E\lambda; 0_g^+ T_0 T_0 \to \lambda_i^{\pi} T_0+1, T_0) = \frac{1}{2T_0+2} \sum_{\mu} \left| \langle | \Omega_{\lambda\mu i} [T^{(+)}, \Gamma(E\lambda\mu)] | \rangle \right|^2 , \tag{15}$$

where

$$\Gamma(E\lambda\mu) = \sum_{j_1 m_1 j_2 m_2} \langle j_2 m_2 | \tilde{\Gamma}(E\lambda\mu) | j_1 m_1 \rangle a^{+}_{j_2 m_2} a_{j_1 m_1} ,$$

$$\tilde{\Gamma}(E\lambda\mu) = i^{\lambda} e \left\{ (\tfrac{1}{2} - t_z) e_p^{(\lambda)} + (\tfrac{1}{2} + t_z) e_n^{(\lambda)} \right\} z^{\lambda} Y_{\lambda\mu}(\theta, \varphi) .$$

Here $e_n^{(\lambda)}$ and $e_p^{(\lambda)}$ are effective charges for neutron and proton. After transformations for $e_p^{(\lambda)} - e_n^{(\lambda)} = 1$ we get

$$B(E\lambda; 0_g^+ T_0 T_0 \to \lambda_i^{\pi} T_0+1, T_0) =$$

$$= \frac{e^2}{2} \frac{1}{2T_0+2} \left| \sum_{j_p j_n} f^{(\lambda)}(j_p j_n) (v_{j_p} u_{j_n} \psi^{\lambda i}_{j_p j_n} + u_{j_p} v_{j_n} \varphi^{\lambda i}_{j_p j_n}) \right|^2 . \tag{16}$$

Similarly we can derive an expression for the reduced $M\lambda$-transition probabilities.

3. Characteristics of the giant dipole resonances

Now we calculate the B(E1)-values for the $T_>$ giant dipole-resonances (GDR) in several spherical nuclei. For comparison we present the results of calculations for $T_<$ GDR. The giant isovector dipole $T_<$ resonance has been calculated in ref./9/ within the RPA. Rather a good description of the energy centroids of the GDR has been obtained. To corectly describe its width, one should take into account the quasiparticle-phonon interactions. Such calculations have been performed in refs. /9,10/, and a sufficiently good description of the form of the GDR has been obtained.

In our calculations we use the energies and wave functions of the Saxon-Woods potential. The potential parameters and the pairing constants are taken from ref./11/. The values of the chemical potentials for the neutron and proton systems for the states $T_>$ of the GDR are the same as for the ground states of doubly even nuclei. For these states the number of particles is conserved with an accuracy higher than 90%. The excited $T_>$ state energy is

$$\mathcal{E}_i = \omega_{\lambda i} + \Delta E_c . \tag{17}$$

where the Coulomb energy ΔE_c is calculated by

$$\Delta E_c = 1.444 (Z - \tfrac{1}{2}) / A^{1/3} - 1.131 \quad \text{MeV} . \tag{18}$$

The results of the RPA calculations of the B(E1, ↑ i)-values for the GDR in ^{88}Sr, ^{90}Zr and ^{92}Mo are shown in fig. 1 and in ^{116,120,124}Sn in fig. 2. It is seen from fig. 1 that the $T_<$ 1^- state strength is concentrated in the region of 14-18 MeV with maximum at 17 MeV, that is in agreement with

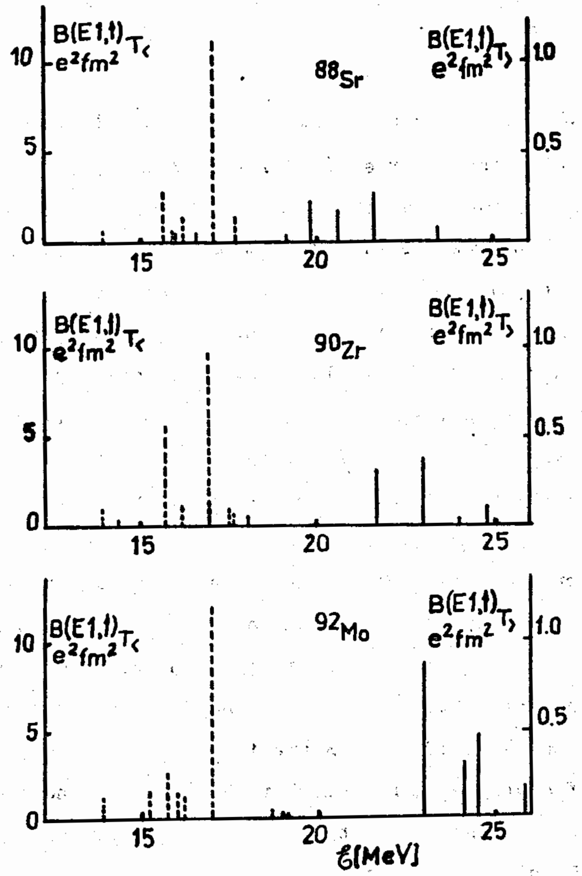

Fig. 1. Reduced E1-transition probabilities for excitation of $T_<$ and $T_>$ components of the GDR in ^{88}Sr, ^{90}Zr and ^{92}Mo.

Notation: ---- is $T_<$ components (the scale is to the left), ——— is $T_>$ components (the scale is to the right).

the experimental data of refs./24-27/. The $T_>$ states lie at 20-23 MeV in ^{88}Sr, 21-25 MeV in ^{90}Zr and 23-26 MeV in ^{92}Mo. According to the experimental data/25-28/ the $T_>$ GDR in these nuclei is at 20-21 MeV. In the Sn isotopes the GDR has been

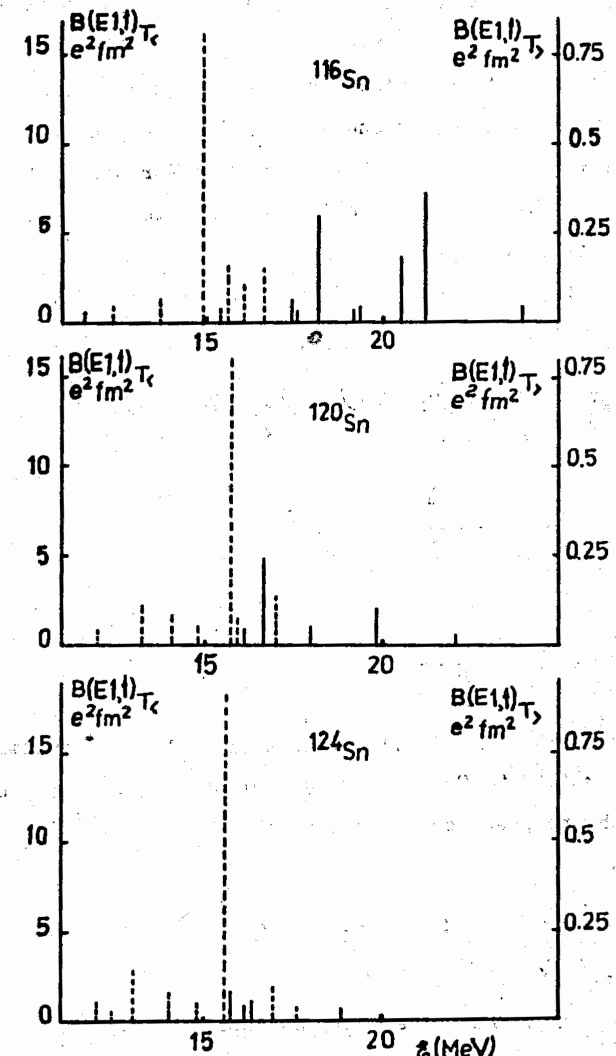

Fig. 2. Reduced E1-transition probabilities for excitation of $T_<$ and $T_>$ components of the GDR in $^{116,120,124}Sn$

For notation see fig. 1.

studied experimentally in ref./29-31/. The $T_<$ state strength is concentrated in the region of 12-20 MeV with maximum at 16 MeV. According to our calculations the most collective 1^- $T_<$

state lies at 15-16 MeV. The 1^- $T_>$ states in the Sn isotopes have been studied experimentally in refs./30-32/ and analysed in ref./27/. For the $T_>$ states of GDR in the Sn isotopes the energy somewhat increases and the B(E1)-value decreases considerably with increasing A. The tendency for the B(E1)-values to decrease with increasing A is seen from the results of our calculations (fig. 2).

The integral characteristics of the $T_>$ GDR as the energy centroids

$$\bar{E}_> = \sum_i B(E1,T_>,i)\mathcal{E}_i \Big/ \sum_i B(E1,T_>,i) \tag{19}$$

and ratio

$$\frac{\sigma_{-1}(T_>)}{\sigma_{-1}(T_<)}, \tag{20}$$

where

$$\sigma_{-1} = \sum_i \frac{\sigma(i)}{\mathcal{E}_i}$$

are shown in the table. It presents also the experimental data from ref./27/. The results of our calculations are compared with the phenomenological estimates of ref./33/, according to which

$$\frac{\sigma_{-1}(T_>)}{\sigma_{-1}(T_<)} = \frac{1}{T_0}\,\frac{1-1.5\,T_0A^{-2/3}}{1+1.5A^{-2/3}}. \tag{21}$$

It is seen from the table that the energy centroids $\bar{E}_>$ are well described, though in the Sn isotopes $\bar{E}_>$ decreases with increasing A, that is inconsistent with the experimental data. The calculated ratio $\sigma_{-1}(T_>)/\sigma_{-1}(T_<)$ agrees with the experimental data for ^{88}Sr, ^{90}Zr, ^{92}Mo and three times as large as the experimental values for ^{116,120,124}Sn. Our calculated ratio $\sigma_{-1}(T_>)/\sigma_{-1}(T_<)$ is $1^1/_2$-2 times as less as the calculations

Table. Integral characteristics $T_>$ of GDR

Nucleus	$\bar{E}_>$ (MeV)		$\sigma_{-1}(T_>)/\sigma_{-1}(T_<)$		
	exp.	calc.	exp.	calc.	calc.*)
^{88}Sr	21.5	20.6	0.06	0.04	0.084
^{90}Zr	20.5	22.2	0.009	0.08	0.117
^{92}Mo	21.0	23.6	0.11	0.10	0.164
^{116}Sn	20.0	19.9	0.009	0.04	0.058
^{120}Sn	20.5	17.8	0.007	0.02	0.036
^{124}Sn	21.0	17.4	0.003	0.01	0.022

*) calculation by formula (21)

by formula (21). It decreases four times when passing from ^{116}Sn to ^{124}Sn, whereas for the experimental data it decreases three times. This ratio decreases more rapidly due to the factor $(T_0+1)^{-1}$ and it is also caused by the decreases in the coefficients U^i_{jn}. The decrease in the cross sections $\sigma(T_>)$ in the isotopes - with increasing has been explained theoretically in ref./34/.

We can state that a good description of the $T_<$ and $T_>$ states of the GDR is obtained with the same values of the isovector constant of the dipole-dipole interaction.

Thus, in this paper we have shown that within the quasiparticle-phonon nuclear model one can calculate the characteristics of the $T_>$ giant resonances. The model apparatus can easily be generalized for calculating the fragmentation of $T_>$ one-phonon states and thus for describing in more detail the strength distribution of $T_>$ components of the giant multipole and spin-multipole resonances.

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