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# FRAGMENTATION OF THE GAMOW-TELLER RESONANCE IN SPHERICAL NUCLEI

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## 1. INTRODUCTION

The progress has been achieved in the experimental study of the giant charge-exchange resonances [1-5]. The Gamow-Teller(GT) and charge-exchange $\Delta\ell = 1$ resonances are observed in many nuclei. At the early stage the investigation of the GT resonance was related with the calculation of the $\beta$-decay strength functions [6-8]. In recent years the energies of charge-exchange resonances have been calculated in the RPA in spherical nuclei with closed shells [9-13].

A part of strength of the charge-exchange resonances is observed experimentally as compared to the relevant sum rules. The quenching of the GT and charge-exchange $\Delta\ell = 1$ transition strength has two reasons, namely the strength fragmentation due to the coupling with the 2p-2h states and to the mixing with the $\Delta$-isobar - nucleon hole configuration. Fiegig and Wambach [14] have calculated the GT resonance on $^{208}$Pb taking into account the admixtures of the $\Delta$-hole and 2p-2h configurations. It was found that the $\Delta$-hole admixtures decrease the nucleon GT strength by 27%, i.e., they did not succeed in explaining the missing strength due to the $\Delta$-isobar-hole configuration. The influence of the 2p-2h configurations turned out to be essential, though only the first $2^+_1$ and $3^-_1$ vibrational states in $^{208}$Pb, were taken into account. Adashi [15], Muto et al. [16] have used the Tamm-Dancoff approximation to calculate the GT spreading width on $^{208}$Pb including the 2p-2h states. Bertsch and Hamamoto [17] have investigated the effect of the central and tensor forces on the fragmentation of charge-exchange phonons and obtained that half of the GT strength in $^{90}$Zr is shifted into the energy region of (15-50) MeV. According to Knüpfer and Matsch [18] the inclusion of 2p-2h states with tensor forces in $^{90}$Zr shifts 26% of the strength into the high-energy region. Bortignon et al. [19] studied the influence of 2p-2h configurations on the distribution on the GT strength on $^{40}$Ca. Many authors [20-25] assume that missing of GT strength is due to the admixtures of the $\Delta$-hole excitations to the proton-particle - neutron-hole GT states. Brown and Rho [22] state that a consistent description of the position of the GT resonance in $^{208}$Pb and the renormalization of the strength through coupling to $\Delta$-isobar region, can be achieved in terms of $\pi$- and $\rho$-meson exchange forces. However, the quenching of the GT strength cannot be thus explained.

It is evident that the $\Delta$-hole configurations play an important role in quenching the strength of charge-exchange phonons. Nevertheless, one should study more thoroughly the fragmentation of GT phonons to clarify how the fragmentation is responsible for the missing of the GT strength in the maximum region. The calculations should not be restricted to magic nuclei only, as it has been made before. One should perform calculations with pairing but without tensor forces for some spherical nuclei. The solution of this problem is the aim of the present paper. In further investigations one can additionally take into account tensor forces and the coupling with the $\Delta$-hole configurations.

The fragmentation of charge-exchange phonons is studied within the quasiparticle-phonon nuclear model /26/. In this case the neutron-proton **np** phonon operators are used, which have been introduced by Kuzmin and Soloviev /27/ for describing the $T_>$ part of the giant dipole resonance. A general method for introducing **np** or charge-exchange phonons in the quasiparticle-phonon model is presented by Soloviev /28/. In the same paper the author obtained the general equations for the fragmentation of charge-exchange one-phonon states in spherical nuclei and the system of approximate equations used for a description of the fragmentation of charge-exchange phonons. The present paper contains the basic approximate equations of the model, which allow one to describe charge-exchange resonances. In Sec.2 we introduce the charge-exchange phonons and obtain the RPA secular equation. Section 3 is devoted to the description of the fragmentation of charge-exchange phonons within the quasiparticle-phonon nuclear model and to the calculation of strength of the (p, n) transitions. The obtained results are discussed in Sec. 4, and the conclusions are made in Sec.5.

## 2. CHARGE-EXCHANGE PHONONS AND RPA EQUATIONS

The Hamiltonian of the quasiparticle-phonon nuclear model consists of the average field as the Saxon-Woods potential for neutron and proton, superconducting pairing interactions and the effective separable residual interaction. To generate the state of a natural parity, we use the multipole-multipole forces

$$V^{M}_{\lambda\mu} = -[\kappa^{\lambda}_{0} + \kappa^{\lambda}_{1}(\vec{\tau}_1,\vec{\tau}_2)]R_{\lambda}(r_1)R_{\lambda}(r_2)Y^{+}_{\lambda,-\mu}(\theta_1,\phi_1)Y_{\lambda,-\mu}(\theta_2,\phi_2). \qquad (1)$$

The spin-multipole - spin-multipole interaction is chosen in the form

$$V^{SM}_{L\lambda\mu} = -[\kappa^{L\lambda}_{0} + \kappa^{L\lambda}_{1}(\vec{\tau}_1,\vec{\tau}_2)]R_{L\lambda}(r_1)R_{L\lambda}(r_2)[Y_L(\theta_1,\phi_1)\sigma_1]^{+}_{\lambda,-\mu}[Y_L(\theta_2,\phi_2)\sigma_2]_{\lambda,-\mu}, \qquad (2)$$



where $[Y_L(\theta,\phi)\sigma]_{\lambda\mu} = \sum_{M,m}\langle LM1m|\lambda\mu\rangle V_{LM}(\theta,\phi)\sigma^m$. The spin-multipole forces are responsible for the states of unnatural parity at $L=\lambda\pm1$ and natural parity at $L=\lambda$. Here we shall use the forces with $L=\lambda-1$ only. The radial dependence of residual forces may be arbitrary. Usually $R_\lambda(r)=r^\lambda$ or $R_\lambda(r)=\partial U(r)/\partial r$ ($U(r)$ is the central part of the average field potential). The constants: isoscalar $\kappa^{\lambda}_{0}$, $\kappa^{L\lambda}_{0}$ and isovector $\kappa^{\lambda}_{1}$, $\kappa^{L\lambda}_{1}$ are determined from either the experimental data or phenomenological estimates. Isotopic matrices $(\vec{\tau}_1,\vec{\tau}_2)$ can be separated into the charge-neutral and charge-exchange interaction parts. The charge-neutral interaction is used for calculating the fragmentation of phonon forming the giant electric and magnetic resonances /26,29-30/. The charge-exchange interaction can be written as follows:

$$H^{ch}_{res} = -2\sum_{\lambda\mu}\kappa^{\lambda}_{1}\beta^{+}_{\lambda\mu}\beta_{\lambda\mu} - 2\sum_{\lambda,\mu}\kappa^{\lambda-1\lambda}_{1}\beta^{+}_{\lambda-1,\lambda\mu}\beta_{\lambda-1,\lambda,\mu}, \qquad (3)$$

where

$$\beta_{\lambda\mu}(\lambda-1,\lambda,\mu) = \sum_{j_n m_n, j_p m_p}\langle j_p m_p|O_{\lambda\mu}(\lambda-1,\lambda,\mu)\,t^{(-)}|j_n m_n\rangle a^{+}_{j_p m_p}a_{j_n m_n},$$

$$O_{\lambda\mu} = i^{\lambda}r^{\lambda}Y_{\lambda,-\mu}(\theta,\phi),\quad O_{\lambda-1)\lambda,\mu} = i^{\lambda-1}r^{\lambda-1}[Y_{\lambda-1}(\theta,\phi)\sigma]_{\lambda,-\mu},$$

$a_{jm}$ ($a^{+}_{jm}$) are the nucleon annihilation (creation) operators and $j_p m_p$ ($j_n m_n$) are quantum numbers of the proton (neutron) single-particle states.

Now we perform the canonical Bogolubov transformation $a_{jm} = u_j\alpha_{jm} + (-)^{j-m}v_j\alpha^{+}_{j,-m}$ and introduce the phonon operators

$$\Omega_{\lambda\mu i} = \sum_{j_n j_p}(\psi^{\lambda i}_{j_p j_n}A(j_p j_n;\lambda\mu) - (-)^{\lambda-\mu}\phi^{\lambda i}_{j_p j_n}A^{+}(j_p j_n;\lambda-\mu)), \qquad (4)$$

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

In the RPA the ground state wave function of a doubly even nucleus is approximated by the phonon vacuum $|\rangle$

$$\Omega_{\lambda\mu i}|\rangle = 0 \qquad (5)$$

and the phonon operators satisfy the commutation relations

$$\langle|[\Omega_{\lambda\mu i},\Omega^{+}_{\lambda'\mu' i'}]_{-}|\rangle = \delta_{\lambda\lambda'}\delta_{\mu\mu'}\sum_{j_p j_n}(\psi^{\lambda i}_{j_p j_n}\psi^{\lambda i'}_{j_p j_n} - \phi^{\lambda i}_{j_p j_n}\phi^{\lambda i'}_{j_p j_n}) = \delta_{\lambda\lambda'}\delta_{\mu\mu'}\delta_{ii'}. \qquad (6)$$

Using (4) and (6), we get

$$\beta_{\lambda-1,\lambda,\mu} = \sum_{j_p j_n}\frac{f(\lambda;j_p j_n)}{\sqrt{2\lambda+1}}\{(-)^{\lambda-\mu}u_{j_n}u_{j_p}B^{+}(j_p j_n;\lambda\mu) - v_{j_n}v_{j_p}B(j_p j_n;\lambda-\mu)\} +$$

$$+\frac{1}{2\sqrt{2\lambda+1}}\sum_i \{D^{(+)}_{\lambda i}[(-)^{\lambda-\mu}\Omega_{\lambda\mu i}+\Omega^+_{\lambda,-\mu i}]+D^{(-)}_{\lambda i}[(-)^{\lambda-\mu}\Omega_{\lambda\mu i}-\Omega^+_{\lambda,-\mu i}]\}, \quad (7)$$

where

$$D^{(+)}_{\lambda i}=\sum_{j_pj_n} f(\lambda;j_pj_n)u^{(+)}_{j_pj_n}g^{\lambda i}_{j_pj_n}, \qquad D^{(-)}_{\lambda i}=\sum_{j_pj_n} f(\lambda;j_pj_n)u^{(-)}_{j_pj_n}w^{\lambda i}_{j_pj_n},$$

$$g^{\lambda i}_{j_pj_n}=\psi^{\lambda i}_{j_pj_n}+\phi^{\lambda i}_{j_pj_n}, \quad w^{\lambda i}_{j_pj_n}=\psi^{\lambda i}_{j_pj_n}-\phi^{\lambda i}_{j_pj_n}, \quad u^{(\pm)}_{j_pj_n}=u_{j_p}v_{j_n}\pm v_{j_p}u_{j_n},$$

$$B(j_pj_n:\lambda\mu)=\sum_{m_pm_n}(-)^{j_n+m_n}\langle j_pm_pj_nm_n|\lambda\mu\rangle a^+_{j_pm_p}a_{j_n,-m_n},$$

and $f(\lambda;j_pj_n)=\langle j_p||O_{\lambda-1,\lambda}t^{(-)}||j_n\rangle$ is the reduced matrix element of the spin-multipole operator. Formulae for the multipole phonons can be derived by changing the matrix elements and effective constants.

Like in the book by Soloviev [33], to find the RPA equations, we use the variational principle

$$\delta\{\langle|\Omega_{\lambda\mu i}H\Omega^+_{\lambda\mu i}|\rangle-\langle|H|\rangle-\frac{\tilde\omega_{\lambda i}}{2}[\sum_{j_pj_n}g^{\lambda i}_{j_pj_n}w^{\lambda i}_{j_pj_n}-1]\}=0 \quad (8)$$

and get

$$g^{\lambda i}_{j_pj_n}=\frac{2}{2\lambda+1}\kappa_1^{\lambda-1,\lambda}f(\lambda;j_pj_n)\frac{\epsilon(j_pj_n)u^{(+)}_{j_pj_n}D^{(+)}_{\lambda i}+\tilde\omega_{\lambda i}u^{(-)}_{j_pj_n}D^{(-)}_{\lambda i}}{\epsilon^2(j_pj_n)-\tilde\omega^2_{\lambda i}},$$

$$w^{\lambda i}_{j_pj_n}=\frac{2}{2\lambda+1}\kappa_1^{\lambda-1,\lambda}f(\lambda;j_pj_n)\frac{\tilde\omega_{\lambda i}u^{(+)}_{j_pj_n}D^{(+)}_{\lambda i}+\epsilon(j_pj_n)u^{(-)}_{j_pj_n}D^{(-)}_{\lambda i}}{\epsilon^2(j_pj_n)-\tilde\omega^2_{\lambda i}},$$

where $\epsilon(j_pj_n)=\epsilon(j_p)+\epsilon(j_n)$ and $\epsilon(j)$ is the one-quasiparticle energy. Inserting expressions for $g^{\lambda i}_{j_pj_n}$ and $w^{\lambda i}_{j_pj_n}$ into the definitions $D^{(\pm)}_{\lambda i}$, we get the system of equations, which has non-zero solutions only at $\tilde\omega_{\lambda i}$ satisfying the equation

$$\mathcal{F}(\tilde\omega)\equiv(\kappa_1^{\lambda-1,\lambda}X^{(+)}_\lambda(\tilde\omega)-1)(\kappa_1^{\lambda-1,\lambda}X^{(-)}_\lambda(\tilde\omega)-1)-(\kappa_1^{\lambda-1,\lambda}X^{(0)}_\lambda(\tilde\omega))^2=0,$$

$$X^{(\pm)}_\lambda(\tilde\omega)=\frac{2}{2\lambda+1}\sum_{j_pj_n}\frac{[f(\lambda,j_pj_n)u^{(\pm)}_{j_pj_n}]^2\epsilon(j_pj_n)}{\epsilon^2(j_pj_n)-\tilde\omega^2}, \quad (9)$$

$$X^{(0)}_\lambda(\tilde\omega)=\frac{2\tilde\omega}{2\lambda+1}\sum_{j_pj_n}\frac{f^2(\lambda,j_pj_n)u^{(+)}_{j_pj_n}u^{(-)}_{j_pj_n}}{\epsilon^2(j_pj_n)-\tilde\omega^2}.$$



The phonon amplitudes are

$$\psi^{\lambda i}_{j_pj_n}=\frac{1}{\sqrt{2\lambda+1}}\sqrt{\frac{\kappa^{\lambda-1,\lambda}}{\mathcal{Y}_{\lambda i}}}\frac{u^{(+)}_{j_pj_n}+\tilde y_{\lambda i}u^{(-)}_{j_pj_n}}{\epsilon(j_pj_n)-\tilde\omega_{\lambda i}}f(\lambda;j_pj_n),$$

$$\phi^{\lambda i}_{j_pj_n}=\frac{1}{\sqrt{2\lambda+1}}\sqrt{\frac{\kappa_1^{\lambda-1,\lambda}}{\mathcal{Y}_{\lambda i}}}\frac{u^{(+)}_{j_pj_n}-\tilde y_{\lambda i}u^{(-)}_{j_pj_n}}{\epsilon(j_pj_n)+\tilde\omega_{\lambda i}}f(\lambda;j_pj_n), \quad (10)$$

$$\mathcal{Y}_{\lambda i}=\frac{\frac{\partial}{\partial\tilde\omega}\mathcal{F}(\tilde\omega)|_{\omega=\tilde\omega_{\lambda i}}}{\kappa_1^{\lambda-1,\lambda}X^{(-)}_{\lambda i}(\tilde\omega_{\lambda i})-1}, \quad \tilde y_{\lambda i}=-\frac{\kappa_1^{\lambda-1,\lambda}X^{(0)}_\lambda(\tilde\omega_{\lambda i})}{\kappa_1^{\lambda-1,\lambda}X^{(-)}_\lambda(\tilde\omega_{\lambda i})-1}.$$

Consider the single particle operators

$$O^{(-)}_{\lambda-1,\lambda,\mu}=\sum_{j_pm_pj_nm_n}\langle j_pm_p|O_{\lambda-1,\lambda,-\mu}t^{(-)}|j_nm_n\rangle a^+_{j_pm_p}a_{j_nm_n} \quad (11)$$

and

$$O^{(+)}_{\lambda-1,\lambda,\mu}=\sum_{j_pm_pj_nm_n}(-)^{\lambda-\mu}\langle j_nm_n|O_{\lambda-1,\lambda,\mu}t^{(+)}|j_pm_p\rangle a^+_{j_nm_n}a_{j_pm_p}. \quad (11')$$

The operator $O^{(-)}_{\lambda-1,\lambda,\mu}$ ($O^{(+)}_{\lambda-1,\lambda,\mu}$) describes single-particle transitions from the target-nucleus $N_0, Z_0$ to the nuclear states $N_0-1$, $Z_0+1$ ($N_0+1$, $Z_0-1$) with decreasing (increasing) by unity isospin projection. The amplitudes of these transitions are

$$\Phi^{(\pm)}_{\lambda i}\equiv\sqrt{2\lambda+1}\langle|\Omega_{\lambda\mu i}Q^{(\pm)}_{\lambda'-1,\lambda',\mu'}|\rangle=$$
$$=\delta_{\lambda\lambda'}\delta_{\mu\mu'}\frac{1}{2}\sqrt{2\lambda+1}\frac{1\pm\tilde y_{\lambda i}}{\sqrt{\kappa_1^{\lambda-1,\lambda}\mathcal{Y}_{\lambda i}}}. \quad (12)$$

The one-phonon states

$$\Omega^+_{\lambda\mu i}|\rangle \quad (13)$$

are a superposition of excitations in nuclei with $N_0-1$, $Z_0+1$ and $N_0+1$, $Z_0-1$. In each state one can separate a dominating branch. Mixing is due to the pairing and ground states correlations. If there is no pairing in the neutron or proton system, eq. (9) splits.

Taking into account the secular equations for the one-phonon states, the Hamiltonian of the quasiparticle-phonon nuclear model is

$$H_M = \sum_{jm} \epsilon(j) a^+_{jm} a_{jm} + H_{mv} + H_{mvq} + H_{smu} + H_{smvq} + \tag{14}$$
$$+ H_{cm} + H_{csmv} + H_{cmvq} + H_{csmvq} \; ,$$

the explicit form of $H_{mv}$, $H_{smv}$, $H_{mvq}$ and $H_{smvq}$ is given in [26,34]

$$H_{csmv} = - \sum_{\lambda\mu i i'} \frac{1 - \bar{y}_{\lambda i} \bar{y}_{\lambda i'}}{\sqrt{\mathcal{Y}_{\lambda i} \mathcal{Y}_{\lambda i'}}} \Omega^+_{\lambda\mu i} \Omega_{\lambda\mu i'} , \tag{15}$$

$$H_{csmvq} = - \sum_{\lambda\mu i} \sqrt{\frac{\kappa^{\lambda-1,\lambda}}{\mathcal{Y}_{\lambda i}}} \sum_{j_p j_n} \frac{f(\lambda; j_p j_n)}{\sqrt{2\lambda+1}} \{ (-)^{\lambda-\mu} u_j u_j B(j_p j_n : \lambda\mu) - \tag{16}$$
$$- v_{j_p} v_{j_n} B^+(j_p j_n; \lambda-\mu) \} \{ (1 - \bar{y}_{\lambda i}) \Omega^+_{\lambda,-\mu i} + (1 + \bar{y}_{\lambda i}) (-)^{\lambda-\mu} \Omega_{\lambda\mu i} \} + \text{h.c.}$$

The terms $H_{cmv}$ and $H_{cmvq}$ differ from (15) and (16) by the matrix elements and isovector constants. The model Hamiltonian is constructed so that all operators $A(jj'; \lambda\mu)$, $A^+(jj'; \lambda\mu)$, $A(j_p j_n; \lambda\mu)$ and $A^+(j_p j_n: \lambda\mu)$ are expressed through the phonon operators, and the quasiparticle operators enter only in the form of $B(jj': \lambda\mu)$, $B(j_p j_n: \lambda\mu)$ and $B^+(j_p j_n: \lambda\mu)$. Under such a construction of the Hamiltonian there is no double counting. In nuclear field theory a special procedure developed to overcome double counting has been developed by Bes et al. [35].

## 3. FRAGMENTATION OF CHARGE-EXCHANGE PHONONS

To study the fragmentation of charge-exchange phonons, we write down the wave function of an odd-odd nucleus as

$$|JM\nu\rangle = \Big[ \sum_i R_i^{J\nu} \Omega^+_{JMi} + \sum_{\lambda_1 i_1, \lambda_2 i_2} P^{\lambda_1 i_1}_{\lambda_2 i_2}(J\nu) [Q^+_{\lambda_1 i_1} \Omega^+_{\lambda_2 i_2}]_{JM} \Big] |\rangle , \tag{17}$$

where

$$[Q^+_{\lambda_1 i_1} \Omega^+_{\lambda_2 i_2}]_{JM} = \sum_{\mu_1, \mu_2} \langle \lambda_1 \mu_1 \lambda_2 \mu_2 | JM \rangle Q^+_{\lambda_1 \mu_1 i_1} \Omega^+_{\lambda_2 \mu_2 i_2}$$

and the phonon creation operator

$$Q^+_{\lambda_1 \mu_1 i_1} = \frac{1}{2} \sum_{\tau = n,p} \sum_{jj'}{}^{\tau} (\psi^{\lambda i}_{jj'} A(jj': \lambda\mu) - (-)^{\lambda-\mu} \phi^{\lambda i}_{jj'} A^+(jj': \lambda-\mu))$$

as in [26,29-32].

The general form of equations for describing the fragmentation is given in paper by Soloviev [28]. In this paper we use the boson commutators for the phonon operators and the condition $[Q_{\lambda_1 \mu_1 i_1}, \Omega^+_{\lambda_2 \mu_2 i_2}]_- = 0$. Then the normalization condi-



tion of the wave function (17) is

$$\xi_\nu \equiv \sum_i (R_i^{J\nu})^2 + \sum_{\lambda_1 i_1, \lambda_2 i_2} (P^{\lambda_1 i_1}_{\lambda_2 i_2}(J\nu))^2 = 1 . \tag{18}$$

The coherent transition strength to the state (17) has the form

$$B^{(\pm)}(J\nu) \equiv \sum_{M_1, M_2} |\langle JM_1 \nu | O^{(\pm)}_{JM_2} | \rangle|^2 = \sum_{ii'} R_i^{J\nu} R_{i'}^{J\nu} \Phi^{(\pm)}_{Ji} \Phi^{(\pm)}_{Ji'} . \tag{19}$$

The energies $\eta_\nu$ and coefficients $R_i^{J\nu}$ and $P^{\lambda_1 i_1}_{\lambda_2 i_2}(J\nu)$ are determined by the variational principle

$$\delta \{ \langle JM\nu | H_M | JM\nu \rangle - \eta_\nu (\xi_\nu - 1) \} = 0 , \tag{20}$$

where the Hamiltonian $H_M$ is taken in the form of (14).

As a result we get the system of linear homogeneous equations

$$\sum_{i'=1}^{N} a_{ii'}(\eta_\nu) R_{i'}^{J\nu} = 0 , \tag{21}$$

where

$$a_{ii'}(\eta) = (\tilde{\omega}_{Ji} - \eta) \delta_{ii'} - \sum_{\lambda_1 i_1, \lambda_2 i_2} \frac{V^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji) \, V^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji')}{\omega_{\lambda_1 i_1} + \tilde{\omega}_{\lambda_2 i_2} - \eta} , \tag{22}$$

N is the number of one-phonon components in the wave function (17), and

$$V^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji) = \langle | \Omega_{JMi} H_M [Q^+_{\lambda_1 i_1} \Omega^+_{\lambda_2 i_2}]_{JM} | \rangle .$$

The system of equations (21) has nonzero solutions only for $\eta_\nu$:

$$\tilde{\mathfrak{F}}(\eta_\nu) \equiv \det || a_{ii'}(\eta_\nu) || = 0 . \tag{23}$$

One can easily show that

$$R_i^{J\nu} = \frac{(-)^{i+i'} M_{ii'}(\eta_\nu)}{M_{ii}(\eta_\nu)} \Big[ - \frac{M_{ii}(\eta_\nu)}{(\partial/\partial\eta) \tilde{\mathfrak{F}}(\eta) |_{\eta=\eta_\nu}} \Big]^{1/2} , \tag{24}$$

where $M_{ii'}$ are the minors of the determinant (23).

Thus, the coherent transition strength to the states (17) is

$$B^{(\pm)}(J, \nu) = - \frac{\sum_{i,i'=1}^{N} (-)^{i+i'} M_{ii'}(\eta_\nu) \Phi^{(\pm)}_{Ji} \Phi^{(\pm)}_{Ji'}}{(\partial/\partial\eta) \tilde{\mathfrak{F}}(\eta) |_{\eta=\eta_\nu}} . \tag{25}$$

The following strength function was calculated to simplify the numerical calculations:

$$b^{(\pm)}(J,\eta) = \sum_{\nu} \rho_{\Delta}(\eta-\eta_{\nu}) B^{(\pm)}(J,\nu), \quad \rho_{\Delta} = \frac{\Delta}{2\pi}\,\frac{1}{(\eta-\eta_{\nu})^2+\Delta^2/4} . \tag{26}$$

The averaging energy interval $\Delta$ determines the way of presentation of the results of calculation. Using the analytical properties of (25) and the weight function $\rho_{\Delta}(\eta-\eta_{\nu})$ we have

$$b^{(\pm)}(J,\eta) = \frac{1}{\pi}\,\mathrm{Im}\,\frac{\sum\limits_{j,j'=1}^{N}(-)^{j+j'} M_{jj'}(\eta+i\frac{\Delta}{2})\,\Phi^{(\pm)}_{Jj}\,\Phi^{(\pm)}_{Jj'}}{\tilde{\mathcal{F}}(\eta+i\frac{\Delta}{2})} . \tag{27}$$

The obtained formulae can be used to describe the fragmentation of any charge-exchange phonons; and the strength functions, to calculate the (p, n) and (n, p) transition strength.

## 4. RESULTS AND DISCUSSION

The parameters of the Saxon-Woods potential and the pairing constants are the same as in[30]. Other parameters of the Saxon-Woods potential were also used for the calculation of $^{90}Zr$ and $^{208}Pb$. The isovector Gamow-Teller interaction constant $\kappa_1^{01}$ (see formula (3)) is equal to -23/A MeV (see[3]). The change of the constant $\kappa_1^{01}$ from - 20/A MeV to - 25/A MeV causes a small change of the RPA phonon energies and a certain redistribution of the transition strength. The multipole and spin-multipole ordinary and neutron-proton phonons with $\lambda^{\pi} = 1^{\pm}$, $2^{\pm}$,... $7^{\pm}$, $8^{-}$ were used as phonons in the two-phonon terms of the wave function (17). From several dozen up to several hundred roots of the RPA secular equations were used for each $\lambda^{\pi}$. The isoscalar and isovector constants of the multipole-multipole and spin-multipole - spin-multipole interactions were either fixed from the experimental data or calculated by the phenomenological formulae like in[30,31]. The present calculations used 10-15 one-phonon terms and up to 2000 two-phonon terms in the wave function (17). The density of two-phonon poles increases strongly with excitation energy. When calculating the strength functions (26) and (27) the averaging energy interval $\Delta = 0.5$ MeV.

Using the strength functions (27) we introduce the total strengths of (n, p) and (p, n) transitions

$$S_{\pm} = \int_0^{E_m} b^{(\pm)}(J,\eta)\,d\eta . \tag{28}$$

The Gamow-Teller transitions have a model independent sum rule in the form

$$S_{-}^{\ell=0} - S_{+}^{\ell=0} = 3(N-Z) . \tag{29}$$



According to our calculations, for the spherical nuclei from $^{90}Zr$ to $^{208}Pb$ the ratio $S_{+}^{\ell=0}/S_{-}^{\ell=0}$ = 0.02-0.03. According to the RPA calculations, under integration in (28) up to 30 MeV $S_{-}^{\ell=0}$ are equal to (96-98)% of the sum-rule limit $3(N-Z)$. These values agree with the values of $S_{-}^{\ell=0}$ for $^{90}Zr$ and $^{208}Pb$ calculated in[13] with a continuum.

As in[29] the Pauli principle is roughly taken into account in the two-phonon terms of the wave function (17). The terms (17) consisting of two noncollective terms, in which the Pauli principle is violated, are neglected. Thus, in calculating $V^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji)$, the terms violating the Pauli principle are excluded. It has been shown[32] that an approximate inclusion of the Pauli principle in calculating the giant electric resonances turned out to be almost equivalent to the exact inclusion of the Pauli principle in the two-phonon components of the wave functions.

The fragmentation of the Gamow-Teller resonance is calculated for several spherical nuclei including nuclei with open shells. The results of calculations by formula (27) (denoted by $\Omega^{+} + Q^{+}\Omega^{+}$) of the (p, n) transition strength functions with excitation of the $1^{+}$ states in odd-odd nuclei and the results of the RPA calculations are given in the table and figs. 1-7. Excitation energies are relative to parent doubly even ground states. In the distribution of the Gamow-Teller strength one can separate the low-energy part, region of maximum and the high-energy part. The table presents the distribution of the Gamow-Teller strength over these regions, which has been calculated in the RPA and by formula (27) taking into account the quasiparticle-phonon interaction. The results are given as per cent of the sum-rule limit $3(N-Z)$. According to the RPA calculations 56-81% of $3(N-Z)$ is concentrated in the region of the resonance maximum. The calculation of the fragmentation of charge-exchange phonons $1^{+}$ leads to the decrease in strength up to 33-60% in the region of the resonance maximum. The values given in the table for $S_{-}^{\ell=0}$ are equal to (90-98)% of $3(N-Z)$, the decrease of the values to (90-92)% is due to the elimination of weak RPA states in the calculations with the quasiparticle-phonon interaction.

Now we discuss the strength distribution of the Gamow-Teller resonance in some nuclei. According to the experimental data in $^{90}Zr$ the energy of the GT resonance maximum is 15.6 MeV and the centroid energy of the low-energy part is 9.2 MeV. Our RPA calculations (see the table and figs. 1 and 2) and calculations in[11-13,23,25] give a right position of the maximum and low-energy part of the GT resonance, in which (70-85)% and (10-20)% of strength is concentrated, respectively. The calculations of $b^{(-)}(1^{+},\eta)$ by formula (27), taking into account the fragmentation of (n, p) phonons, result in the

Table

Strength distribution of the Gamow-Teller resonance

| Nucleus | Energy region of maximum of the GT resonance MeV | Approximation | Strength distribution in % of sum-rule limit 3(N-Z) | | | $S_-^{(-)}$ in % of 3(N-Z) |
|---|---|---|---|---|---|---|
| | | | Below the region of maximum of the GT resonance | Region of maximum of the GT resonance | Above the region of maximum of the GT resonance | |
| $^{90}Zr$ | 12-18 | RPA | 14 | 81 | 3 | 98 |
| | | $\Omega^+ + Q^+\Omega^+$ | 15 | 54 | 29 | |
| $^{120}Sn$ | 12-18 | RPA | 27 | 56 | 7 | 90 |
| | | $\Omega^+ + Q^+\Omega^+$ | 29 | 36 | 25 | |
| $^{124}Sn$ | 14-20 | RPA | 18 | 67 | 5 | 90 |
| | | $\Omega^+ + Q^+\Omega^+$ | 26 | 33 | 31 | |
| $^{124}Te$ | 13-19 | RPA | 25 | 62 | 5 | 92 |
| | | $\Omega^+ + Q^+\Omega^+$ | 31 | 40 | 21 | |
| $^{140}Ce$ | 15-20 | RPA | 12 | 7 | 4 | 94 |
| | | $\Omega^+ + Q^+\Omega^+$ | 30 | 46 | 18 | |
| $^{208}Pb$ | 17-22 | RPA | 11 | 78 | 3 | 92 |
| | | $\Omega^+ + Q^+\Omega^+$ | 19 | 59 | 14 | |

change of distribution of the GT strength. Notably, in $^{90}Zr$ 26% of the GT strength is shifted into the high-energy region and only 54% of strength remains in the region of maximum. In the low-energy region the GT strength is shifted towards lower excitation energies. It should be noted that for a more correct description of the strength distribution by formula (27) in the low-energy region one should impose a condition for the strength not to be lower than the ground state energy of a doubly odd nucleus.

Bertch and Hamamoto /17/ investigated the GT strength distribution in $^{90}Zr$ taking into account central and tensor forces. They considered a large number of the 2p-2h states and used a perturbative treatment. They found that roughly half of the GT strength is removed from the resonance towards higher excitation energies up to 50 MeV. The tensor forces are



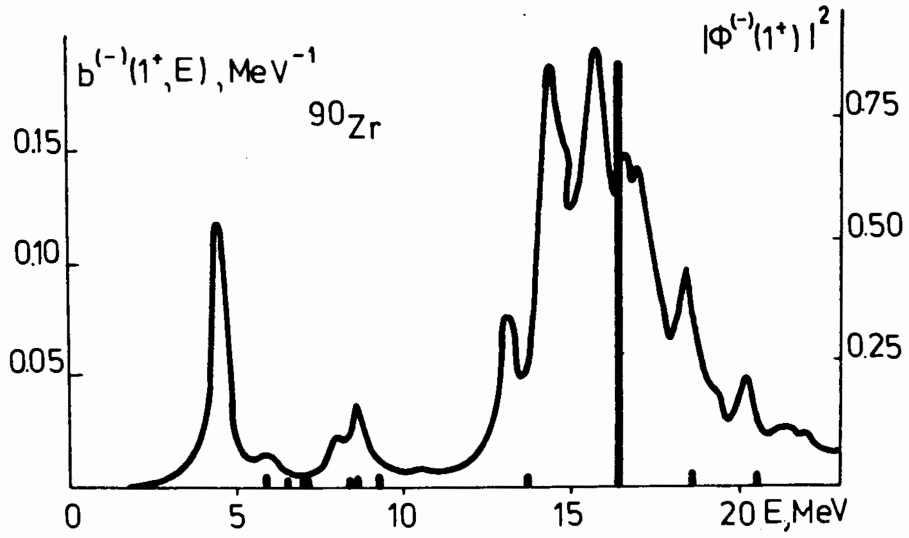


Fig.1. Fragmentation of the Gamow-Teller resonance on $^{90}Zr$. Parameters of the single-particle potential are from paper by Pyatov and Fayans /13/. Vertical lines are the results of the RPA calculations. Solid curve is the strength function including two-phonon admixtures. Both the scalars are normalized to 3(N-Z). Energies are reckoned from the ground $^{90}Zr$ state.

responsible for about 1/3 of the effect obtained. According to the more consistent calculations of Knüpfer and Metsch /18/ in $^{90}Zr$ 74% of the GT strength is concentrated in the region of maximum and only 26% is shifted into the high excitation energies. In the calculations of Knüpfer and Metsch /18/ and in our calculation almost the same part of the GT strength is shifted towards the high-energy region. In our calculations the tensor forces are neglected but a large number of one-phonon states is used, and the GT strength is mainly shifted towards the region of (18-25) MeV. According to the papers /17, 18/ the GT strength is distributed up to 50 MeV. It is possible that the tensor forces are responsible for the shift of a small part of the GT strength up to 50 MeV. The analysis performed by Scholten et al. /38/ indicates the presence of the GT strength at high excitation energies.

The details of distribution of the Gamow-Teller strength on $^{90}Zr$ depend on the position of single-particle levels. As is known (see, for instance, /37/ ), there are some difficulties in describing the low-lying and deep hole states in $^{89}Zr$ and $^{91}Mo$. It is difficult to choose the parameters so as to describe the position of levels in these nuclei. The dependence of the fragmentation of the GT resonance on the position of single-particle levels is demonstrated in figs. 1 and 2. The calculations with the parameters of the Saxon-Woods potential of Pyatov and Fayans [13], are shown in fig. 1, whereas with the parameters [30] in fig.2. The comparison of figs.1 and 2 shows that the use of the parameters [30] causes a splitting of the GT resonance and a larger shift of the low-energy part of strength towards lower energies. The fragmentation of one-phonon states depends on the completeness of the one-phonon basis. Figure 2 shows the calculations in which the two-phonon terms (17) include phonons with $\lambda^{\pi} = 1^{\pm}, 2^{\pm}, \ldots 7^{\pm}, 8^{-}$ and phonons with $\lambda^{\pi} = 1^{\pm}, 2^{\pm}, 3^{\pm}$. The comparison of the results of calculations with the complete and reduced one-phonon basis shows that the inclusion of phonons with $\lambda \geq 3$ increases considerably the fragmentation of one-phonon states. Possibly, the role of phonons with $\lambda > 3$ in our calculations is somewhat overestimated.

The distribution of the GT strength in $^{120,124}$Sn, $^{124}$Te and $^{140}$Ce is shown in the table and figs.3-6. Our RPA calculations

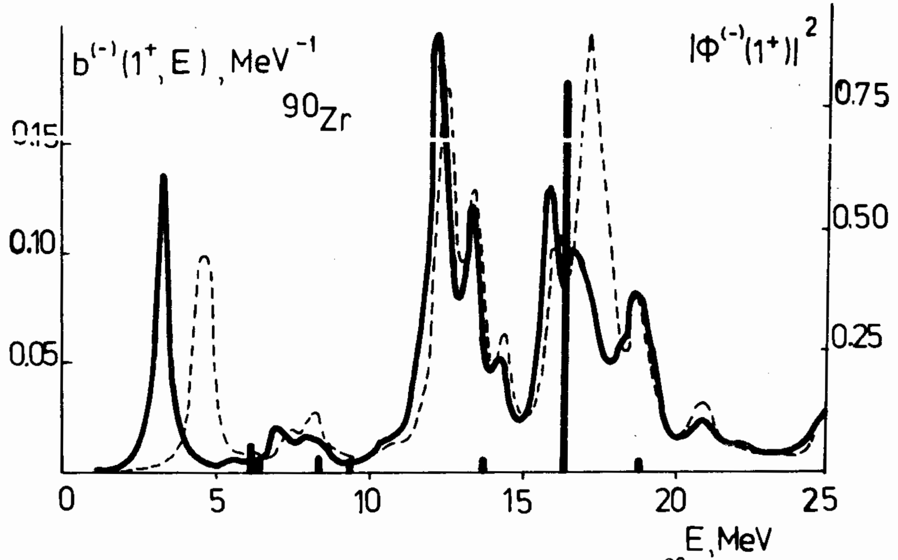


Fig.2. Fragmentation of the GT resonance on $^{90}$Zr. Parameters of the single-particle potential are from paper by Ponomarev et al. [30]. Solid line is the strength function of GT transitions (in the two-phonon part (17) the phonons with $\lambda^{\pi} = 1^{\pm}, 2^{\pm}, \ldots 7^{\pm}, 8^{-}$ are included). Dashed line is the strength function of GT transitions (in the two-phonon part the phonons with $\lambda^{\pi} = 1^{\pm}, 2^{\pm}, 3^{\pm}$ are included). The rest notation is as in fig.1.



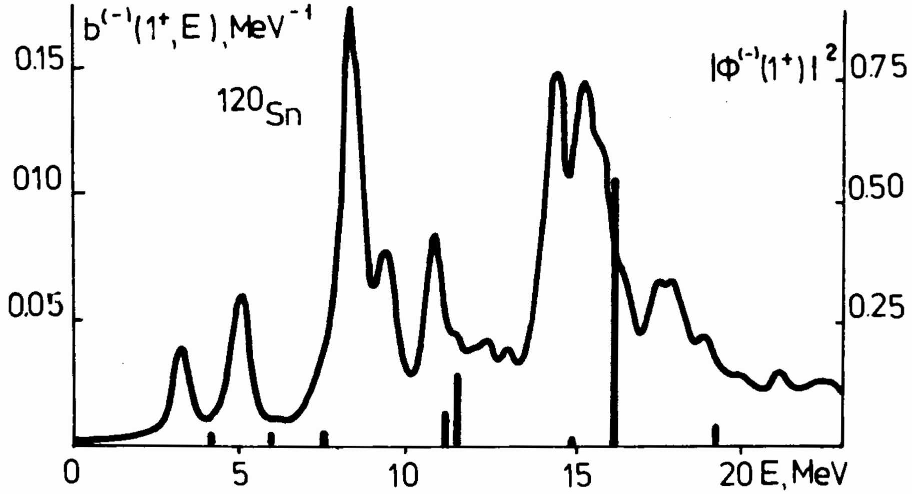


Fig.3. Fragmentation of the GT resonance on $^{120}$Sn. The notation is as in fig.1.

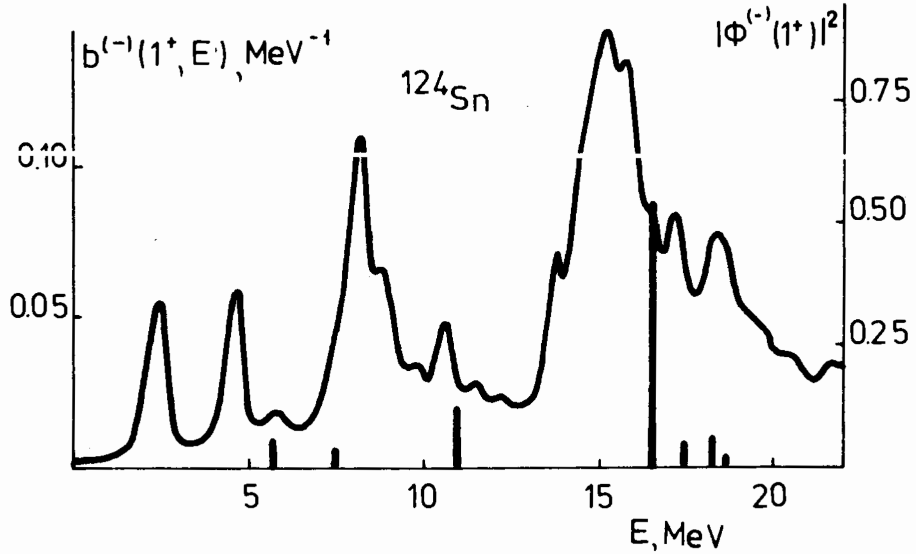


Fig.4. Fragmentation of the GT resonance on $^{124}$Sn. The notation is as fig.1.

for $^{120}$Sn are in agreement with the calculations of Sagawa and Nguyen van Giai [23], according to which the strength in the low-energy region amounts to about a half of the GT strength in the region of maximum. Due to the quasiparticle-phonon interaction the GT strength decreases down to 36% in the region of resonance maximum, increases up to 25% in the high-energy region and is shifted towards lower excitation energies in the

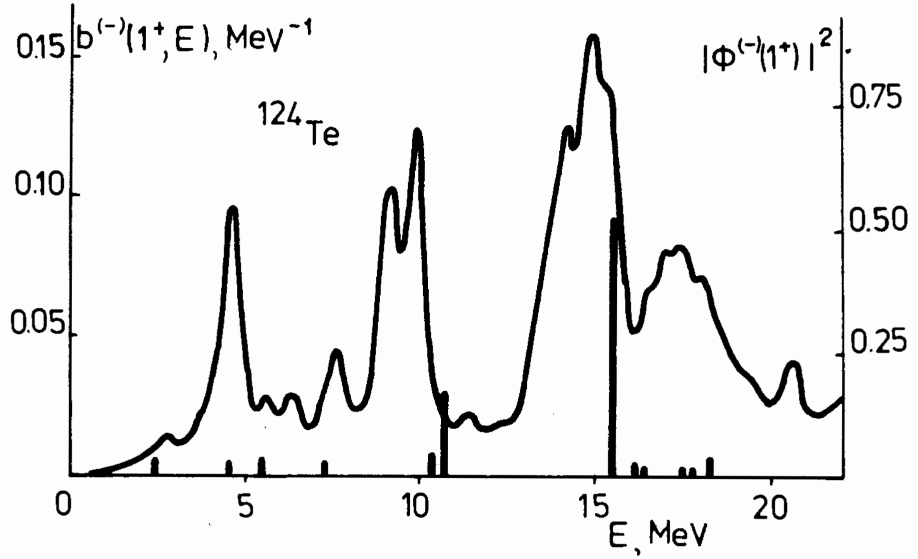


Fig.5. Fragmentation of the GT resonance on $^{124}Te$. The notation is as in fig.1.

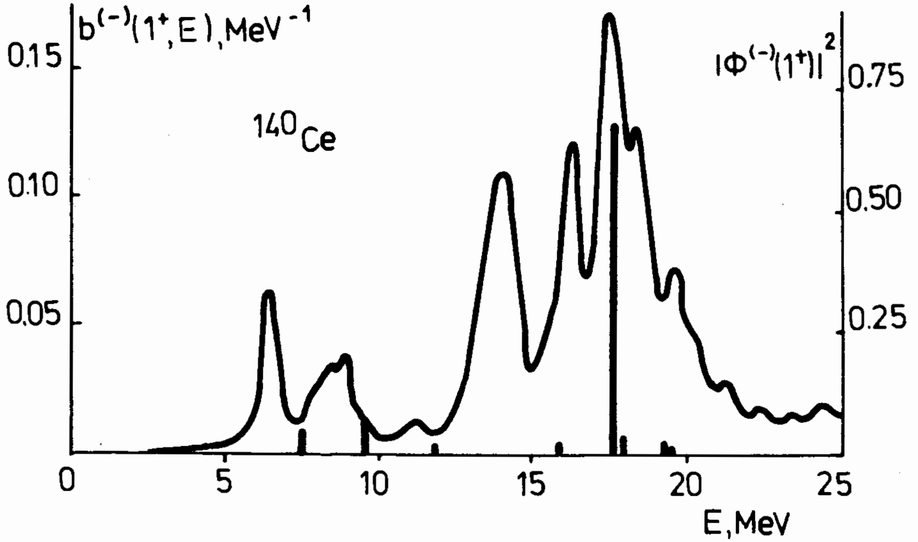


Fig.6. Fragmentation of the GT resonance on $^{140}Ce$. The notation is as in fig.1.



low-energy region. A similar picture is observed in the distribution of the GT strength in $^{124}Sn$. The quasiparticle-phonon interaction shifts 8% of strength from the region of resonance maximum to the low-energy region. The GT strength in the high-energy region. increases up to 31%.

The highly excited states of spherical nuclei with open shells are calculated within the quasiparticle-phonon nuclear model. Such an example is the distribution of the GT strength in $^{124}Te$. Due to the quasiparticle-phonon interaction, the GT strength decreases from 62% to 40% of 3(N - Z) in the region of resonance maximum. In a narrower region of 13-16 MeV the GT strength descreases from 58% to 30%, and in the region of (16-20) MeV 20% of the GT strength is concentrated. There is no considerable difference in the fragmentation of one-phonon states in $^{124}Te$ in comparison with $^{120,124}Sn$ and $^{140}Ce$.

Our RPA calculations in $^{140}Ce$ somewhat differ from the calculations of Sagawa and Nguyen van Giai [23], in which the concentration of the GT strength in the low-energy region is twice as small as in the region of resonance maximum. According to our RPA calculations 12% of the GT strength is in the low-energy region. Due to the quasiparticle-phonon interaction, 18% of the GT strength is shifted into the region of 12-15 MeV, 14% of the GT strength is shifted into the region above 20 MeV. The centroid energies in the region of resonance maximum and low-energy region are 17 and 11 MeV, that is in agreement with the experimental values 16.3 and 10.7 MeV measured by Orihara et al. [38]. The ratio of the GT strength in the region of resonance to the low-energy region, which is 1.5, does not contradict the value 1.2 measured by Orihara et al. However, our calculations indicate the distribution of the GT strength over several levels.

The distributions of the GT strength in $^{208}Pb$ are calculated with the single-particle energies and wave functions of the Saxon-Woods potential with the parameters [31] and the parameters fitted by Voronov. Our RPA calculations and the calculations in [13,14,23,25] give a right position of maximum of the GT resonance and differ by the details of distribution of the GT strength in the low-energy region. The strength distribution in the low-energy region is most sensitive to the position of single-particle levels near the Fermi energy. For the parameters of the potential as in [31] 6.6% of 3(N-Z) is concentrated in the low-energy region whereas for modified parameters 11.3% of 3(N - Z).

Due to the quasiparticle-phonon interaction a part of the GT strength is shifted into the low-energy and high-energy region. According to the calculations with the parameters of the Saxon-Woods potential [31], the maximum of the GT resonance

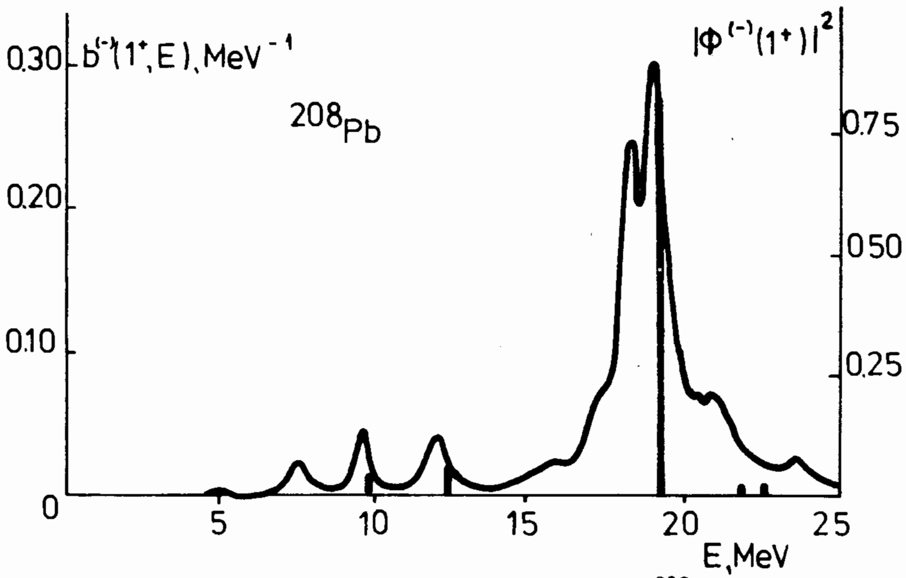


Fig.7. Fragmentation of the GT resonance on $^{208}$Pb. The notation is as in fig.1.

is not splitted. The results of calculations with modified parameters of the Saxon-Woods potential are shown in the table and fig.7. From the region of (17-22) MeV, 11% of the GT strength is shifted to the high-energy region. In the region (17-20) MeV, 49.6% of the GT strength is concentrated, 21.4% of the GT strength is shifted from this region towards higher energies. According to our calculations the GT strength in the region of maximum is larger than the experimental data. An additional missing of the GT strength is obviously due to the coupling with the Δ isobar-nucleon hole configurations.

Consider now the distribution of the GT strength in $^{208}$Pb we have obtained with the calculations including the 2p-2h configurations. According to our calculations the distribution of the GT strength in the region of maximum has no such a fine structure as in the calculations of Bortignon et al./39/, and it is in a qualitative agreement with the data of Muto et al./16/. The fragmentation of neutron-proton one-phonon states is not so strong as that obtained by Adashi/15/. The distribution of the GT strength we obtained differs from that obtained by Wambach and Schwesinger/40/ by form in the low-energy part and by value in the energy region higher than 25 MeV. The quasiparticle-phonon interaction leads almost to the same distribution of the GT strength in the low-energy region as the effect treated by Grotz et al./41/.



Thus, we may conclude that within the quasiparticle-phonon nuclear model the distribution of the GT strength in doubly magic nuclei $^{90}$Zr and $^{208}$Pb is described not worse than within other models.

## CONCLUSION

The quenching of the Gamow-Teller strength in the region of its maximum is caused by two mechanisms. They are the fragmentation of one-phonon neutron-proton states due to the coupling with more complex configurations and to the mixing with the isobar - nucleon hole configuration. In calculating the fragmentation of the GT one-phonon states due to the coupling with two-phonon states the ground state correlations are taken into account. The ground state correlations are not taken into account in the calculations including the coupling with 2p-2h configurations. Usually, the central forces are taken into account in the calculations. The calculations performed in this paper have shown that the quasiparticle-phonon interactions are responsible for the shift of (10-30)% of strength from the region of maximum towards larger excitation energies. As a result, (20-30)% of the GT strength appeared in the energy interval (20-30) MeV. In this way the authors succeeded to explain a considerable part of the quenching of the GT strength. It is possible that the tensor forces being included cause a shift of a certain part of the GT strength into the region of higher excitation energies up to 50 MeV. The missing of another part of the GT strength is obviously due to the mixing with the Δ-isobar - nucleon hole configuration.

The investigation of the fragmentation of neutron-proton phonons within the quasiparticle-phonon nuclear model will be continued. First, we shall calculate the fragmentation of charge-exchange spin-dipole and E1 states in spherical nuclei. The charge-exchange resonances in deformed nuclei have been calculated by Soloviev et al./42/ in the RPA. The characteristic features of the strength distribution of the GT and charge-exchange spin-dipole and E1 resonances in deformed nuclei are discussed.

In conclusion we should like to note that the quasiparticle-phonon nuclear model serves as a good basis to study the fragmentation of single-particle and collective motions in medium and heavy nuclei.

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Фрагментация гамов-теллеровского резонанса в сферических ядрах

В рамках квазичастично-фононной модели ядра рассчитана фрагментация однофононных состояний, описывающих гамов-теллеровский резонанс, в $^{90}$Zr, $^{120,124}$Sn, $^{124}$Te, $^{140}$Ce и $^{208}$Pb. Сила ГТ резонанса распределена по трем областям: низкоэнергетической, максимума резонанса и высокоэнергетической. Учет взаимодействия квазичастиц с фононами приводит к смещению /20-30/% ГТ силы из области резонанса к более высоким энергиям возбуждения. В результате в области максимума ГТ резонанса остается /40-60/% от величины $3(N-Z)$, следующей из правила сумм. В низкоэнергетической области сосредоточено /20-30/% от $3(N-Z)$. Утверждается, что фрагментация однофононных зарядово-обменных состояний является одной из причин исчезновения ГТ силы в области его максимума.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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Fragmentation of the Gamow-Teller Resonance in Spherical Nuclei

The fragmentation of one-phonon states, describing the Gamow-Teller resonance, in $^{90}$Zr, $^{120,124}$Sn, $^{124}$Te, $^{140}$Ce, and $^{208}$Pb is calculated within the quasiparticle-phonon nuclear model. The GT strength is distributed over three regions: low-energy region, the region of resonance maximum and the high-energy region. The quasiparticle-phonon interaction shifts (10-30)% of the GT strength from the resonance region towards higher excitation energies. As a result, (40-60)% of sum-rule limit $3(N-Z)$ remains in the region of maximum of GT resonance. The low-energy contains (15-30)% of $3(N-Z)$. The fragmentation of one-phonon charge-exchange states is thought to be one of the reasons for the quenching of the GT strength in the region of its maximum.

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

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