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Transversal E1-Resonance in Spherical Nuclei

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Abstract

The spin–flip E1-states in medium and heavy spherical nuclei are investigated within the quasiparticle–phonon nuclear model. The RPA-calculations predict the existence of a collective 1^--state formed by the isovector spin–dipole force. Its excitation energy is about 20 MeV. This state is intensively excited by the inelastic electron scattering at momentum transferred $q \approx 0.5$–$0.7\ \mathrm{fm}^{-1}$, at the scattering angles $\theta > 60°$ the main contribution to its excitation comes from the transversal form factor. Therefore, it can be recognized as the transversal E1-resonance. The interaction with two-phonon states causes a very strong spreading of the resonance, thus making its experimental observation hardly probable.

1. Introduction

The spin and spin–isospin components of the effective forces in nuclei and the relevant nuclear properties are intensively studied at present (see, for instance, [1]). The attenuation is mainly paid to magnetic and charge-exchange resonances and the low-lying excitations. However, according to the selection rules, the spin forces should influence the properties of electric nuclear excitations too. This problem is less studied, though it is known that the influence of the spin forces on the properties of Eλ-states causes interesting phenomena. For instance, the so-called transversal E1 resonance has been observed in the inelastic electron scattering in ^{16}O [2]. This resonance is due to the spin–dipole component of the effective forces [3]. The calculations of Eramzhyan and Goncharova [4] indicate a possible existence of the transversal E1 resonance in heavier nuclei, ^{58,60}Ni, too.

In the present paper we study the properties and the possibility to reveal experimentally the transversal E1 resonance in medium and heavy nuclei. As is known, in these nuclei the properties of high-lying collective excitations are strongly influenced by the interaction with two-phonon (or $2p$–$2h$) configurations. To take this interaction into account, we used the quasiparticle–phonon nuclear model (QPM).

2. Effective interaction

The basic assumptions of the QPM are expounded in papers [5–7]. The model Hamiltonian contains the single-particle potentials for neutrons and protons, pairing forces and separable multipole and spin-multipole forces (isoscalar and isovector). The structure of electric states with momentum and parity $\pi = (-1)^\lambda$ are determined by the following components of effective forces:

$$V(r_1, r_2) = \tfrac{1}{2} \sum_{\mu} (\kappa_0^{(\lambda)} + \kappa_1^{(\lambda)} \tau_1 \tau_2) f(r_1) f(r_2) Y_{\lambda\mu}(\Omega_1) Y^*_{\lambda\mu}(\Omega_2) + (\kappa_0^{(\lambda\lambda)} + \kappa_1^{(\lambda\lambda)} \tau_1\tau_2) f(r_1) f(r_2) [\sigma_1 Y_{\lambda\mu'}(\Omega_1)]_{\lambda\mu} \times [\sigma_2 Y_{\lambda\mu''}(\Omega_2)]^+_{\lambda\mu} \quad (1)$$

where $\kappa_0^{(\lambda)}$, $\kappa_1^{(\lambda)}$ and $\kappa_0^{(\lambda\lambda)}$, $\kappa_1^{(\lambda\lambda)}$ are the isoscalar and isovector constants of multipole and spin-multipole forces. In our calculations the radial form factor $f(r)$ is chosen as a derivative of the central part of the Saxon–Woods potential U: $f(r) = \partial U/\partial r$.

The Hamiltonian parameters are determined from the experimental data on the low-lying collective states and giant resonances of doubly even nuclei. Since in this paper the electric dipole excitations are investigated, we shall dwell upon the choice of the constants $\kappa_{0,1}^{(1)}$ and $\kappa_{0,1}^{(11)}$ only.

The constants of separable dipole forces $\kappa_{0,1}^{(1)}$ are determined by the experimental energy of the giant dipole resonance and under the condition that the energy of the spurious 1^- state vanishes. We have no unambiguous experimental indications for the choice of the $\kappa_{0,1}^{(11)}$ constants. One can use the experimental data on the M2 resonance, since its properties are also determined by the spin–dipole effective interaction. Under the assumption that the constants $\kappa_{0,1}^{(\lambda L)}$ are the same for all λ and L (for $f(r) = \partial U/\partial r$ this assumption is satisfactory at least for small λ and L), one can choose the values of κ_0 and κ_1 so that the data on the energies of M1 and M2 resonances and excitation probabilities of M2 resonance in different nuclei [8, 9] will be described satisfactorily. Almost the same results for magnetic resonances were obtained with a radial form factor of the separable forces $f(r) = r^\lambda$ [10–13]. The detailed comparison of these two types of calculations has been made in [6]. Since the isoscalar spin–dipole interaction slightly influences the integral characteristics of excitations studied, we do not take into consideration this interaction. Under this suggestion the constant $|\kappa_1^{(11)}|$ is 1.5–2 times as large as the constant of isovector dipole forces $|\kappa_1^{(1)}|$.

3. Structure of one-phonon 1^--states

Consider first the properties of 1^--states in the RPA. The one-phonon wave-function is a linear superposition of forward- and backward-going two-quasiparticle components

$$Q^+_{\lambda\mu i}\Psi_0 = \tfrac{1}{2} \sum_{j_1 j_2} \{\psi^{\lambda i}_{j_1 j_2} [\alpha^+_{j_1 m_1} \alpha^+_{j_2 m_2}]_{\lambda\mu} - (-)^{\lambda-\mu} \varphi^{\lambda i}_{j_1 j_2} [\alpha_{j_2 m_2} \alpha_{j_1 m_1}]_{\lambda -\mu}\} \Psi_0 \quad (2)$$

Ψ_0 is the phonon vacuum or the ground state wave function of a doubly even nucleus. The structure of one-phonon 1^--states generated by the effective interaction (1) has been investigated in [14] in detail. Their energies are determined from the condition of equality to zero of the fourth order determinant (if $\kappa_0^{(11)} = 0$, the determinant is of the third order). Several one-phonon 1^--states have a collective structure. Some of them with the largest excitation probabilities $B(\mathrm{E1}, 0^+_{\mathrm{g.s.}} \to 1^-_i)$ are at the energy $E_x = 12$–17 MeV. This is a well-known giant dipole resonance (GDR). The position of the GDR is determined by the isovector dipole constant $\kappa_1^{(1)}$. The isovector spin–

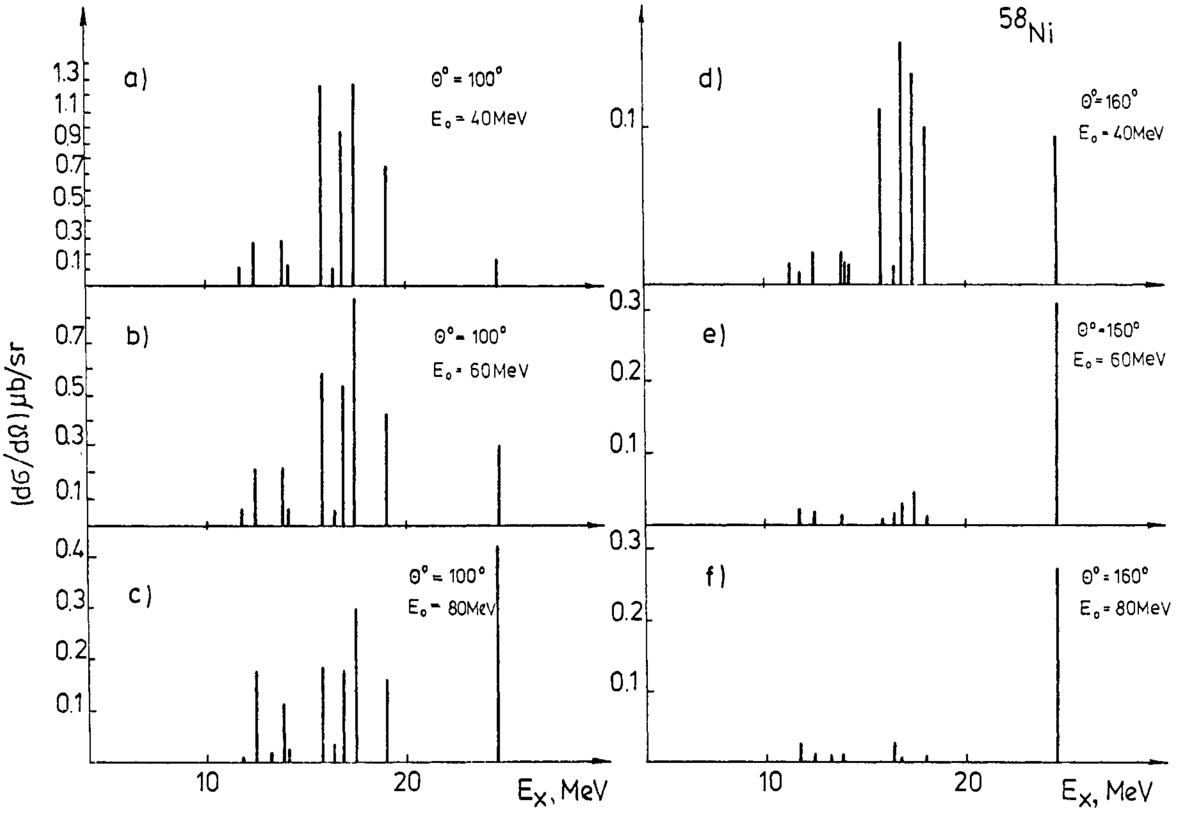


Fig. 1. Electroexcitation cross section of the one-phonon 1^--states in ^{58}Ni for the set of electron energies E_0 and scattering angles θ: (a) $E_0 = 40$ MeV, $\theta = 100°$; (b) $E_0 = 60$ MeV, $\theta = 100°$; (c) $E_0 = 80$ MeV, $\theta = 100°$; (d) $E_0 = 40$ MeV, $\theta = 160°$; (e) $E_0 = 60$ MeV, $\theta = 160°$; (f) $E_0 = 80$ MeV, $\theta = 160°$.

dipole interaction slightly influences its properties changing somewhat B(E1)-distribution over resonance states. A similar result has been obtained earlier in the theory of finite Fermi-systems [15]. Besides GDR-states in all the nuclei studied a strongly collectivized 1^--state appears at higher excitation energy. The maximal contribution of the two-quasiparticle component to its structure does not exceed 25–30%, and its B(E1)-value is less by an order of magnitude than of the GDR-states. The energy of this state decreases with increasing A from $E_x = 24.5$ MeV in ^{58}Ni* up to $E_x = 19.4$ MeV in ^{208}Pb. Thus, the pattern of the E1-strength distribution in the spectrum of a doubly-even nucleus is similar to that shown in Fig. 1(a).

The high-lying collective E1-state appears entirely due to the isovector spin–dipole force. Its structure is formed by the spin–flip single-particle transitions, and when the spin-dipole interaction vanishes ($\kappa_1^{(11)} \to 0$) it disappears as if "breaking up" into several two-quasiparticle states. At the same time the isovector dipole constant $\kappa_1^{(1)}$ slightly influences its properties.

Thus, the isovector spin–dipole interaction may cause the appearance among the E1-excitations of a collective excitation of a new type formed by spin–flip single-particle transitions. The new collective state is by 5–7 MeV higher than the GDR. Both the collective isovector dipole and isovector spin–dipole states are weakly coupled with each other. These properties of dipole and spin-dipole excitations of doubly even nuclei resemble those of charge-exchange 1^--excitations predicted by the theory of finite Fermi-systems [16].

4. Excitation of the spin–flip E1 resonance in the inelastic electron scattering

The collective spin–flip E1 state has a small B(E1)-value, and from this point of view it is not a resonance state like the E1 states forming the GDR. Its resonance properties manifest themselves under certain conditions in the inelastic electron scattering. In light nuclei the spin–flip E1 mode dominates in the (e, e′) scattering cross section at the values of the momentum transferred $q \sim 0.5\ \text{fm}^{-1}$ [3]. Therefore, we shall consider (by the example of ^{58}Ni) the excitation of one-phonon E1-states in the (e, e′)-scattering* at $q = 0.0$–$0.8\ \text{fm}^{-1}$. As a function of the scattering angle θ, this cross section for three one-phonon states with maximal B(E1) (see Fig. 1(a)) and spin–flip collective E1 state is shown in Fig. 2(a). The behaviour of excitation cross sections of well-known resonance E1 states differs from that of the spin–flip E1 state. This is due to essential differences between charge and current transition densities of the relevant excitations. As is known, the charge density of the GDR has a pronounced peak on the nuclear surface. The charge transition density of the spin–flip E1 state $\rho_1(r)$ is concentrated inside the nucleus (see Fig. 3). It is considerably less than the current transition densities $\rho_{10}(r)$ and $\rho_{12}(r)$* determining the behaviour

* The energy of the spin–flip E1-state in ^{58}Ni has been predicted to be equal to 26 MeV in [4]. Note, other effective forces and single-particle spectra have been used.

* The calculations are performed within the DWBA.

* As concerns the definitions of transition densities, we follow paper [17].

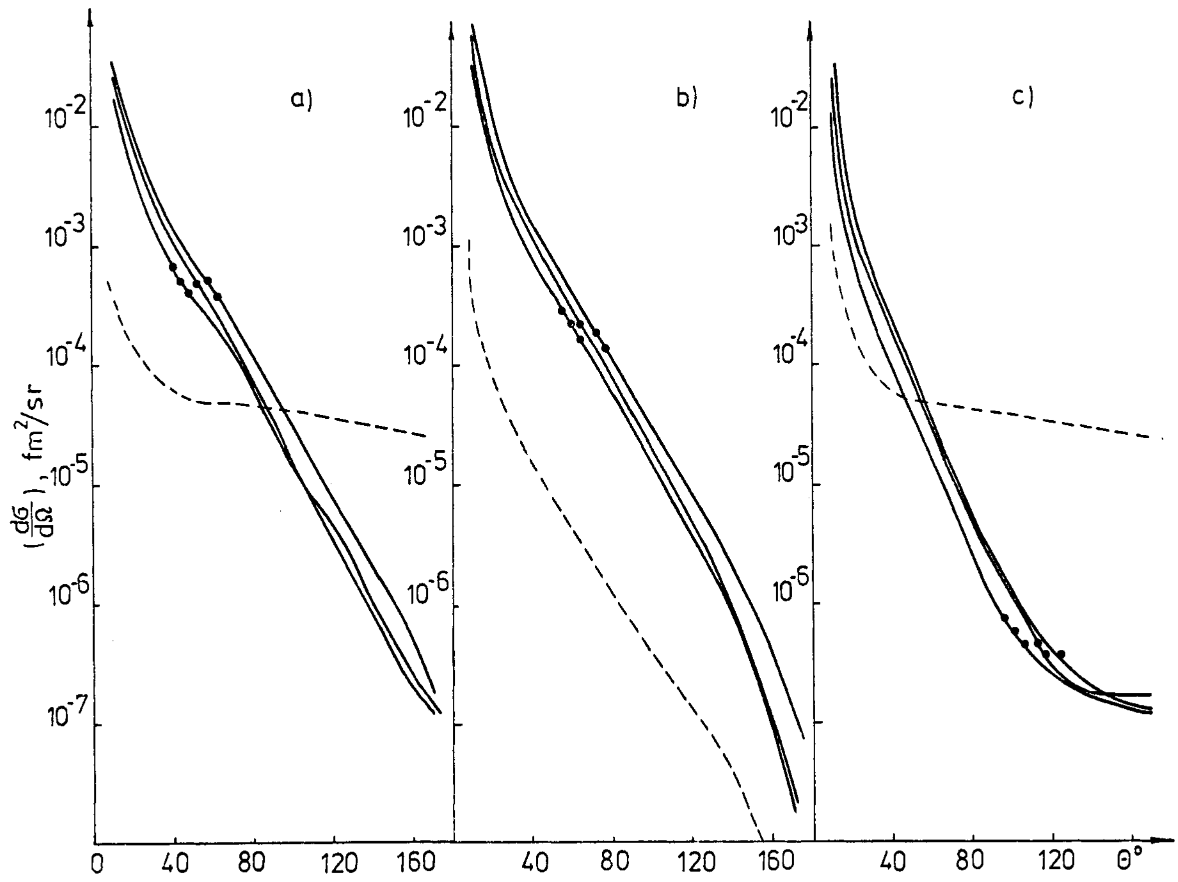


Fig. 2. Electroexcitation cross section (a) and its longitudinal (b) and transversal (c) parts (see eq. (3)) as a function of the scatteing angle θ for the three one-phonon isovector dipole states (full curves) and isovector spin–dipole state (dashed curves) in ^{58}Ni. The energies of the states are $E_x = 15.9$ MeV (•); 17.4 MeV (••); 19.0 MeV (•••) and 24.5 MeV. The electron energy $E_0 = 80$ MeV.

of the transversal form factor. Therefore, for a wide region of scattering angles θ the transversal form factor dominates in the excitation of the spin–flip E1-state. This is clearly seen from Fig. 2. In the Born approximation the (e, e′)-scattering cross section with excitation of the electric state λ^π has the form:

$$(d\sigma/d\Omega)_\lambda \sim V_L(\theta)|F^c_\lambda(q)|^2 + V_T(\theta)|F^T_\lambda(q)|^2 \qquad (3)$$

where $V_L(\theta)$ and $V_T(\theta)$ are the kinematic factors, and $F^c_\lambda(q)$ and $F^T_\lambda(q)$ are the Coulomb and transversal form factors of the state. The contribution to the total cross section of the first and second terms of eq. (3) separately are shown in Fig. 2(b) and (c). The Coulomb form factor of the spin–flip E1-state is less by an order of magnitude than of the GDR-states†. But the dependence of $V_L(\theta)|F^c_{1^-}(q)|^2$ on the angle θ is almost the same for all 1^--states in Fig. 2(b). A quite different situation is for the transversal form factors (Fig. 2(c)). According to the Siegert theorem, at small θ $F^T_{1^-}(q) \sim F^c_{1^-}(q)$ and the behaviour of $F^T_{1^-}(q)$ is the same for all 1^--states too. At $\theta > 60°$ the transversal form factor of the spin–flip state decreases much slower than of the GDR-states. This fact and the relatively small value of $|F^c_{1^-}(q)|^2$ cause an essentially different mechanism of the excitation of the spin– flip state comparing to the GDR-states. The excitation of the latter up to the angles $\theta \sim 150°$ is determined by the Coulomb form factor only. For the isovector spin–flip E1-state the contribution of $F^c_{1^-}(q)$ to the $(d\sigma/d\Omega)_{1^-}$ can be neglected already at $\theta > 60°$.

† This is the reason for such a small B(E1)-value of the spin–flip E1-state.

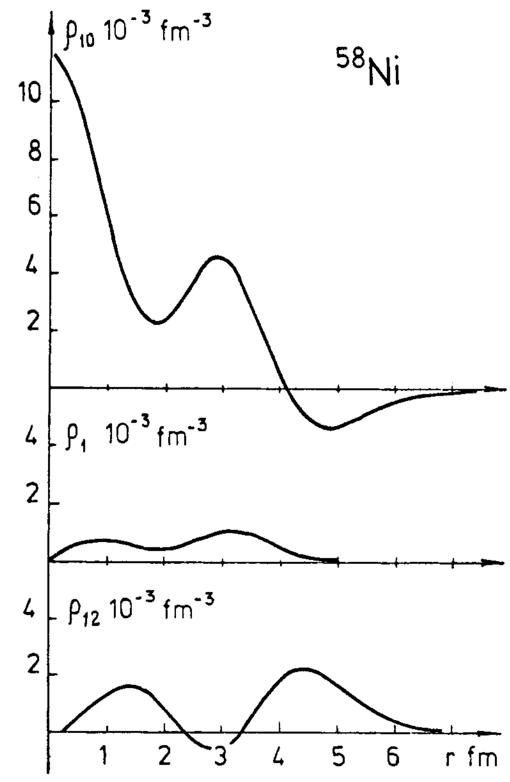


Fig. 3. Current $\rho_{10}(r)$, $\rho_{12}(r)$ and charge $\rho_1(r)$ transition densities of the collective spin–flip E1-state in ^{58}Ni.

Due to the dominant role of the transversal form factor in the excitation of the spin–flip state, it can be recognized as the transversal E1 resonance ($E1_T$ resonance). This term will be used in what follows.

A detailed picture of the excitation probability in the (e, e′)-scattering of one-phonon 1^- states in ^{58}Ni is given in Fig. 1. The excitation cross sections are shown for two scattering angles $\theta = 100°$, $160°$ and for three electron energy values $E_0 = 40$, 60, 80 MeV, that corresponds to the momenta transferred $0.25\ \text{fm}^{-1} \leqslant q \leqslant 0.7\ \text{fm}^{-1}$. Under the choice of θ we have been guided by the following reasons: at $\theta = 100°$ the factor $V_L(\theta)$, which vanishes at $\theta = 180°$, has not yet suppressed the contribution of the Coulomb form factor to the cross section and the GDR excitation cross section has the same order of magnitude as $E1_T$-resonance; at $\theta = 160°$ the values of q correspond to the maximum of $F^T_{1^-}(q)$ already at small E_0. For the smallest value of q (Fig. 1(a)) the distribution of $(d\sigma/d\Omega)_{1^-_{\Gamma}}$ over the nuclear spectrum is almost identical with the distribution of $B(E1,\ 0^+_{g.s.} \to 1^-_i)$ (this is the reason why this figure has been pointed out in Section 3 in discussing the structure of one-phonon 1^- states). In this figure we can clearly see only one resonance region: the region of the GDR ($14 \leqslant E_x \leqslant 19$ MeV). With increasing q (see Fig. 1(b) or 1(d)) the values $(d\sigma/d\Omega)_{1^-}$ decrease, but among many one-phonon 1^--states there appears an intensively excited state at $E_x = 24.5$ MeV. This is just the $E1_T$-resonance. With further increasing q the excitation cross section of the $E1_T$-resonance increases whereas the excitation cross sections of other one-phonon 1^- states continue to decrease. At $q \approx 0.5$–$0.7\ \text{fm}^{-1}$ (Fig. 1(e) and (f)) the transversal E1 resonance dominates in the cross section.

The same pattern is seen in other nuclei too (we have investigated ^{90}Zr, ^{124}Te and ^{208}Pb). We should like to note the following. At $\theta \sim 160°$ and $E_0 \approx 50$–80 MeV the $E1_T$ resonance is excited more intensively than other 1^- states and also than the states with other momenta and parities. This is revealed by our calculations in ^{90}Zr.

Thus, the RPA calculations show that the transversal $E1_T$-resonance may exist in medium and heavy spherical doubly even nuclei as in light nuclei. Its excitation energy varies from 19 to 25 MeV (depending on A). This is a strongly collectivized state, the structure of which is contributed by many particle-hole components corresponding to the spin–flip transitions. It is most appropriate to search for this state experimentally at $q \approx 0.5 \div 0.7\ \text{fm}^{-1}$, that corresponds to the first maximum of the cross section.

5. Quasiparticle–phonon interaction and strength function method

As is known, at large excitation energies the key part belongs to the interaction of one-phonon states with more complex ones, with two-phonon states in the first place. This interaction may drastically change the RPA-results for the distribution of the excitation probability over the nuclear spectrum.

In the QPM the coupling of one- and two-phonon states is caused by the quasiparticle–phonon interaction [7]

$$H_{\text{qph}} = -\tfrac{1}{2} \sum_{\lambda\mu i} \sum_{j_1 j_2} \Gamma(j_1 j_2; \lambda i)\ [Q^+_{\lambda\mu i} - (-)^{\lambda-\mu} Q_{\lambda-\mu i}]\, B(j_1 j_2 \lambda - \mu) + \text{h.c.} \tag{4}$$

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

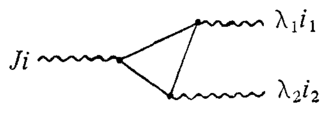
$\Gamma(j_1 j_2;\ \lambda i)$ depends on the structure of phonon $\lambda^\pi i$, reduced matrix element $\langle j_1 \| \partial U/\partial r \| j_2 \rangle$ and the Bogolubov coefficients u_j, v_j. In the wave functions of an excited state one should take into account the admixture of two-phonon components

$$\Psi_\nu(JM) = \left\{ \sum_i R_i(J\nu) Q^+_{JMi} + \sum_{\substack{\lambda_1 i_1 \\ \lambda_2 i_2}} P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu) [Q^+_{\lambda_1 \mu_1 i_1} Q^+_{\lambda_2 \mu_2 i_2}]_{JM} \right\} \Psi_0 \tag{5}$$

The matrix element of the interaction of one- and two-phonon states is

$$\begin{aligned} U^{\lambda_1 i_1}_{\lambda_2 i_2}(Ji) &= \langle \Psi_0 Q_{JMi} \| H_{\text{qph}} \| [Q^+_{\lambda_1 \mu_1 i_1} Q^+_{\lambda_2 \mu_2 i_2}]_{JM} \Psi_0 \rangle \\ &= (-)^{\lambda_1+\lambda_2+J} \frac{1}{\sqrt{2}} [(2\lambda_1+1)(2\lambda_2+1)]^{1/2} \\ &\times \sum_{j_1 j_2 j_3} \left\{ \Gamma(j_1 j_2; Ji) \begin{Bmatrix} \lambda_1 \lambda_2 J \\ j_1 j_2 j_3 \end{Bmatrix} (\psi^{\lambda_1 i_1}_{j_3 j_1} \varphi^{\lambda_2 i_2}_{j_2 j_3} + \psi^{\lambda_2 i_2}_{j_2 j_3} \varphi^{\lambda_1 i_1}_{j_3 j_1}) \right. \\ &+ \Gamma(j_1 j_2; \lambda_1 i_1) \begin{Bmatrix} \lambda_1 \lambda_2 J \\ j_3 j_2 j_1 \end{Bmatrix} (\psi^{Ji}_{j_2 j_3} \psi^{\lambda_2 i_2}_{j_3 j_1} + \varphi^{Ji}_{j_2 j_3} \varphi^{\lambda_2 i_2}_{j_3 j_1}) \\ &+ \left. \Gamma(j_1 j_2; \lambda_2 i_2) \begin{Bmatrix} \lambda_1 \lambda_2 J \\ j_1 j_3 j_2 \end{Bmatrix} (\psi^{Ji}_{j_3 j_1} \psi^{\lambda_1 i_1}_{j_2 j_3} + \varphi^{Ji}_{j_3 j_1} \varphi^{\lambda_1 i_1}_{j_2 j_3} \right\} \end{aligned} \tag{6}$$

Expression (6) is obtained under the usual for the RPA assumption $[Q_{\lambda_1 \mu_1 i_1}, Q^+_{\lambda_2 \mu_2 i_2}] = \delta_{\lambda_1 \lambda_2} \delta_{\mu_1 \mu_2} \delta_{i_1 i_2}$ and corresponds to the following diagram of the nuclear field theory [7, 18, 19]



Corrections to this approximation in the magic and semi-magic nuclei, where 2^+_1 states are not very colletive, are small [20].

To calculate the electroexcitation cross sections of 1^- states eq. (5) we have used the strength function method, suggested in [21] and widely used at present (see [7, 19] and therein). The strength function method for the electroexcitation cross sections is expounded in detail in [11, 22]. We used a standard definition of the strength function:

$$b(d\sigma/d\Omega, E_x) = \sum_\nu \left(\frac{d\sigma}{d\Omega} \right)_{J\nu} \frac{\Delta}{2\pi} \frac{1}{(E_x - \eta_{J\nu})^2 + \Delta^2/4} \tag{7}$$

where $\eta_{J\nu}$ is the energy and $(d\sigma/d\Omega)_{J\nu}$ is the electroexcitation cross section of the state eq. (5), Δ is the parameter. Summation in eq. (7) is performed over all the states eq. (5) with $J^\pi = 1^-$.

6. $E1_T$ resonance breakup as a result of the interaction with two-phonon states

In the two-phonon part of the wave function eq. (5), we took into account phonons with momenta and parities $\lambda^\pi = 1^\pm \div 7^\pm$. Of all possible two-phonon components, allowed by the Pauli principle, we conserved 1000 components with maximal matrix elements $U^{\lambda_2 i_2}_{\lambda_1 i_1}(1^-_i)$. The values $U^{\lambda_2 i_2}_{\lambda_1 i_1}(1^-_i)$ for the rejected components were more than 100 times as small as the maximal value of $U^{\lambda_2 i_2}_{\lambda_1 i_1}(1^-_i)$. The matrix elements and strength functions $b(d\sigma/d\Omega,\ E_x)$ have been calculated by amended version of the GIRES program [23]. However, the terms $\sim \psi\varphi$ and φ^2 in eq. (6) were not taken into account, as always $\psi > \varphi$, and for

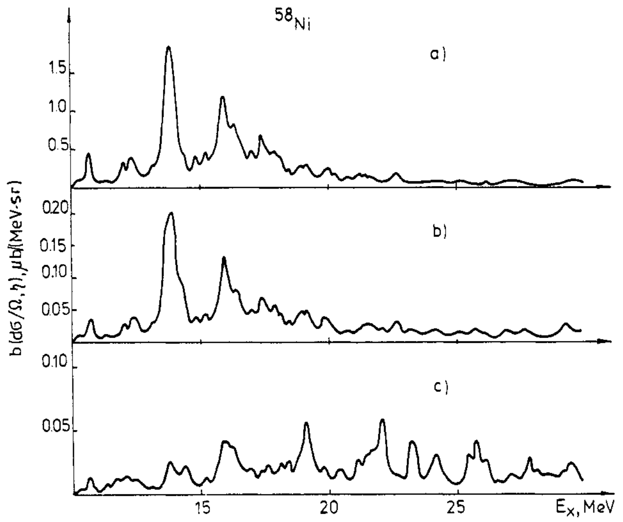


Fig. 4. Electroexcitation strength function of the 1^--states in ^{58}Ni for the set of electron energies E_0 and scattering angles θ: (a) $E_0 = 40$ MeV, $\theta = 100°$; (b) $E_0 = 40$ MeV, $\theta = 160°$; (c) $E_0 = 60$ MeV, $\theta = 160°$.

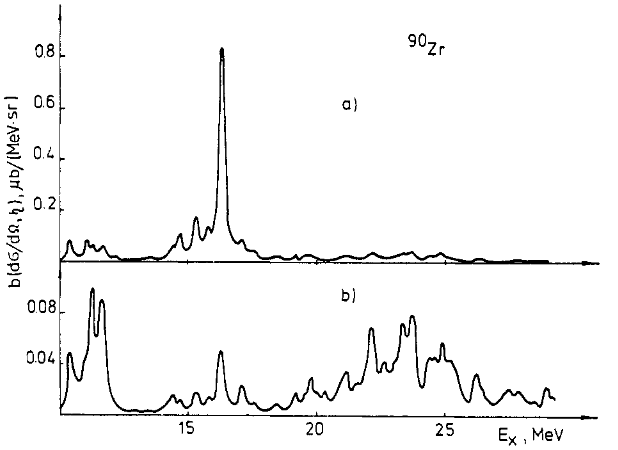


Fig. 5. Electroexcitation strength function of the 1^--states in ^{90}Zr for set of electron energies E_0 and scattering angles θ: (a) $E_0 = 40$ MeV, $\theta = 160°$; (b) $E_0 = 60$ MeV, $\theta = 160°$.

the one-phonon states with an excitation energy of several MeV $\psi \gg \varphi$. It has been tested that large matrix elements $U^{\lambda_2 i_2}_{\lambda_1 i_1}(1^-_i)$ change slightly. In the first sum of eq. (5) both the one-phonon GDR-states and the $\mathrm{E1_T}$ resonance were included. The value of Δ was taken to be equal to 0.5 MeV.

We proceed with the discussion of the results for ^{58}Ni (Fig. 4). The strength function of the electroexcitation cross section eq. (7) in this nucleus has been calculated for three values of q: $E = 40$ MeV, $\theta = 100°$, $160°$; $E_0 = 60$ MeV, $\theta = 160°$. The relevant RPA-results are shown in Fig. 1(a), (d) and (e). As has already been mentioned while discussing the RPA calculations in Section 4, at $E_0 = 40$ MeV and $\theta = 100°$ the $\mathrm{E1_T}$ resonance is excited weakly. Therefore, one can clearly see only one resonance in Fig. 4(a), this is the GDR ($13 \leqslant E_x \leqslant 17$ MeV). The influence of the interaction with two-phonon configurations on this resonance has been discussed many times [18, 24] and we shall not dwell upon this problem. At larger values of q the RPA calculations predicted the appearance of a new resonance, $\mathrm{E1_T}$, in the (e, e′)-scattering cross section (Fig. 1(d)). However, the strength function $b(\mathrm{d}\sigma/\mathrm{d}\Omega, E_x)$ in Fig. 4(b) has no any resonance structure at relevant excitation energies. This means that the one-phonon $\mathrm{E1_T}$ state spread over many states of the type eq. (5). An excitation probability for each this state is very small. The total excitation probability of the $\mathrm{E1_T}$-state is distributed over the energy interval $\Delta E_x \approx 10$ MeV. Since the excitation cross section of the GDR decreases with further increasing q, the strength function in Fig. 4(b) has no any pronounced resonance structures and the cross section is flat.

These results contradict those of [4], where the interaction of the $\mathrm{E1_T}$ resonance with more complex states in ^{58}Ni has been studied too. In these calculations the excitation probability of the $\mathrm{E1_T}$ resonance decreases only three times and the pronounced resonance structure at $E_x \approx 26$ MeV has been preserved. Our calculations show its complete disappearance. The reason for different results lies, as we have already mentioned [11], in a very small number of phonon excitations, the interaction with which has been taken into account in [4]. Therefore, the interaction strength of the $\mathrm{E1_T}$ resonance with two-phonon configurations has been underestimated in [4].

The results for ^{58}Ni are typical. They are equally valid for ^{90}Zr (Fig. 5). At values of q when the giant dipole resonance and $\mathrm{E1_T}$ resonance should be excited with almost the same intensity, the strength function $b(\mathrm{d}\sigma/\mathrm{d}\Omega, E_x)$ has only one pronounced maximum (Fig. 5(a)). This maximum corresponds to the GDR. At values of q, when the GDR form factor has minimum, in the (e, e′)-scattering cross section with excitation of 1^- states there are no any clearly seen resonance structures (Fig. 5(b)).

In conclusion we shall consider the results of calculations for ^{208}Pb. In this nucleus the interaction of one- and two-phonon states, calculated within the QPM, is relatively weak [25]. However, the density of $2p$–$2h$ states at $E_x \sim 20$ MeV is very large and in the one-phonon part of the wave function eq. (5) one should take into account a large number of one-phonon 1^- states. The computational difficulties forced us to include in the wave function eq. (5) only the $\mathrm{E1_T}$ resonance. The giant E1 resonance and $\mathrm{E1_T}$ resonance are slightly coupled with each other. Due to the interaction with two-phonon states their fragmentation proceeds almost independently, that has been tested by us in lighter nuclei. Therefore, an error due to the aforementioned approximation is not large. When in the wave function eq. (5) there is only one one-phonon component, the excitation cross section of this state is proportional to $R^2_\nu(1^-)$ and

$$b(\mathrm{d}\sigma/\mathrm{d}\Omega, E_x) = (\mathrm{d}\sigma/\mathrm{d}\Omega)_{\mathrm{E1_T}}\, b(R^2, E_x)$$

Therefore in Fig. 6 we have depicted the strength function $b(R^2, E_x)$. It is seen that in ^{208}Pb the interaction of the $\mathrm{E1_T}$-resonance with two-phonon states results in its very strong spreading. The excitation probability at maximum decreases several times and becomes comparable with excitation probabilities of other states.

Note, that Fig. 6 is an exciting illustration of how the spreading width of the resonance $\Gamma^\downarrow$ arises due to the interaction with complex configurations. The strength of the only one-phonon $\mathrm{E1_T}$-state is distributed over many levels in the interval $\Delta E_x \sim 5$ MeV.

7. Conclusion

So, our answer to the question whether the transversal E1-resonance does exist in nuclei with $A > 50$ is negative.

Of course, we used a rather schematic effective interaction,

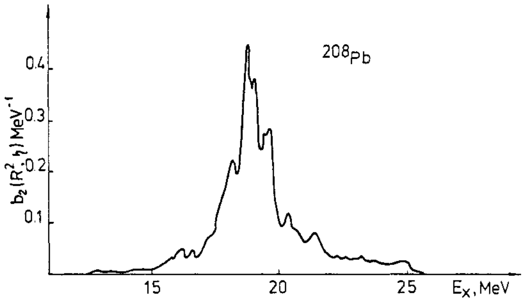


Fig. 6. Strength function of the $E1_T$-resonance in ^{208}Pb.

the parameters of which are not reliably defined in the spin channel. However, it should be noted that the variations of the parameters within rather a large range do not change conclusions made in Sections 3 and 4. The spin–dipole interaction being reduced 2–3 times, the collective $E1_T$-state does not disappear, only its energy decreases by 1–2 MeV. We did not take into account the isoscalar spin–dipole force. There are reasons to believe that it is small [26]. But for $\kappa_0^{(11)} \sim \kappa_1^{(11)}$ as well the basic results of the RPA-calculations are valid. Of course, the separable form of the interaction and its independence of momentum transferred, put the question of a real existence of the collective $E1_T$-state. For example, according to the results of Speth et al. [27] the π- and ρ-exchange leads to such a strong effective attenuation of the spin force that the existence of collective spin excitation in heavy nuclei is hardly probable. We want to stress only that the existence of the $E1_T$-resonance follows from the RPA-calculations with different effective NN-forces. Apart from [4] we should like to mention recent results of Dumitrescu and Suzuki [28] with the SG-II forces.

Our main result, however, consists in that the $E1_T$-resonance disappears due to the coupling with two-phonon states, rather than in the proof of its existence. So, the experimental observation of the $E1_T$-resonance in heavy nuclei is hardly probable. One may imagine as resonances wide structures at $E_x \sim 19$ MeV in ^{208}Pb and $E_x \sim 23$ MeV in ^{90}Zr, but to separate them from the experimental background will be a highly difficult problem. In the only experiment we are aware of where the region of $E_x \sim 20$ MeV in ^{208}Pb has been investigated in the backward (e, e′)-scattering, no pronounced resonance structures have been observed [29].

It should be noted that a complete breakup of the $E1_T$-resonance due to the interaction with two-phonon states is not the only case, the same result has been obtained for the high-lying M2-mode [10, 11]. A strong fragmentation of the Gamow–Teller resonance has been demonstrated by Bertsch and Hamamoto [30]. But the coupling of the high-lying M1-resonance with two-phonon states is not so strong [22].

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References

1. Bertsch, G. F., Nucl. Phys. **A354**, 157c (1981); Ericson, M., Lecture given at the Int. School on Nucl. Physics, Erice, 1983; Ref. TH 3625–CERN (1983).
2. Goldman, A. and Stroetzel, M., Z. Physik **A239**, 235 (1970).
3. Eramzhyan, R. A., Gmitro, M. and Kissener, H. R., Nucl. Phys. **A338**, 436 (1980).
4. Eramzhyan, R. A. and Goncharova, N. G., Z. Physik **A306**, 89 (1982); Goncharova, N. G., Mishchenko, G. M. and Eramzhyan, R. A., Izv. Akad. Nauk SSSR (Ser. Fiz.) **46**, 2091 (1982).
5. Soloviev, V. G., Theory of Complex Nuclei, Oxford: Pergamon, (1976); Soloviev, V. G., Particles and Nuclei **9**, 810 (1978); Soloviev, V. G., Nukleonika, **23**, 1149 (1978).
6. Vdovin, A. I. and Soloviev, V. G., Particles and Nuclei **14**, 237 (1983).
7. Voronov, V. V. and Soloviev, V. G., Particles and Nuclei **14**, 1380 (1983).
8. Richter, A., Proc. Int. School on Nucl. Structure (Alushta, 1980), JINR, D4-80-385, Dubna (1980).
9. Djalali, C., Marty, N., Morlet, M., Willis, A., Jourdain, J. C., Anantaraman, N., Crawley, G. M., Galonsky, A. and Kitching, P., Nucl. Phys. **A388**, 1 (1982).
10. Ponomarev, V. Yu., Soloviev, V. G., Stoyanov, Ch. and Vdovin, A. I., Nucl. Phys. **A323**, 446 (1979).
11. Ponomarev, V. Yu., Stoyanov, Ch. and Vdovin, A. I., J. Phys. G: Nucl. Phys. **8**, L77 (1982).
12. Ponomarev, V. Yu. and Vdovin, A. I., Proc. of the 5th Seminar Electromagn. Interactions of Nuclei at Low and Medium Energies (Moscow, 1981), Moscow, 1982, p. 63; JINR, E4-81-769, Dubna (1981).
13. Castel, B. and Hamamoto, I., Phys. Lett. **65B**, 27 (1976).
14. Nguyen Dinh Dang and Ponomarev, V. Yu., Izv. Akad. Nauk (Ser. Fiz.) **47**, 2147 (1983).
15. Kamerdziev, S. P., Jad. Fiz. **15**, 676 (1972).
16. Pyatov, N. I. and Fayans, S. A., Particles and Nuclei **14**, 953 (1983).
17. Lee, H. C., Chalk River Report, AECL-4839 (1975).
18. Bortingnon, P. F., Broglia, R. A., Bes, D. R. and Liotta, R., Phys. Rep. C **30**, 305 (1977).
19. Bertsch, G. F., Bortignon, P. F. and Broglia, R. A., Rev. Mod. Phys. **55**, 287 (1983).
20. Soloviev, V. G., Stoyanov, Ch. and Nikolaeva, R., JINR, E4-83-166, Dubna (1983); Nguyen Dinh Dang and Ponomarev, V. Yu., Teor. Mat. Fiz. **57**, 154 (1983).
21. Bohr, A. and Mottelson, B., Nuclear Structure, v. 1, W. A. Benjamin, Inc., New York, Amsterdam (1969); Malov, L. A. and Soloviev, V. G., Nucl. Phys. **A270**, 87 (1976).
22. Ponomarev, V. Yu., Stoyanov, Ch. and Vdovin, A. I., Z. Physik **A308**, 157 (1982).
23. Ponomarev, V. Yu., Stoyanova, O. and Stoyanov, Ch., JINR, P4-81-704, Dubna (1981).
24. Soloviev, V. G., Stoyanov, Ch. and Vdovin, A. I., Nucl. Phys. **A288**, 376 (1977); Soloviev, V. G., Stoyanov, Ch. and Voronov, V. V., Nucl. Phys. **A304**, 503 (1978).
25. Soloviev, V. G., Stoyanov, Ch. and Voronov, V. V., Nucl. Phys. **A399**, 141 (1983).
26. Toki, H., Bertsch, G. F. and Cha, D., Phys. Rev. C **28**, 1398 (1983).
27. Speth, J., Klemt, V., Wambach, J. and Brown, G. E., Nucl. Phys. **A343**, 382 (1980).
28. Dumitrescu, T. S. and Suzuki, T., Proc. of the Int. Conf. on Nucl. Structure, v. 1, p. 231. Florence (1983).
29. Woodward, C., Peterson, G. A., Phys. Rev. **C20**, 2437 (1979).
30. Bertsch, G. F., Hamamoto, I., Phys. Rev. **C26**, 1323 (1982).