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Влияние двухфононных примесей на свойства M1 - резонанса в сферических ядрах

Рассчитана силовая функция M1 -переходов в большом числе сферических ядер из интервала массового числа $60 \leq A \leq 140$. Учитывалось взаимодействие одно- и двухфононных состояний. Результаты сравниваются с экспериментом.

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Influence of the Two-Phonon Admixture on the M1 -Resonance in Spherical Nuclei

The influence of the two-phonon admixtures on the M1 -resonance in spherical nuclei with mass numbers $60 \leq A \leq 140$ is studied. The calculations are performed within the quasiparticle-phonon nuclear model with factorized multipole and spin-multipole forces. In nuclei with the number of neutrons 50,82 the role of the two-phonon admixtures is insignificant whereas in other nuclei especially in those with strong pairing in the proton and neutron schemes it is significant. The radiative strength functions are calculated at the neutron binding energy. The results are compared with the experimental data and calculations of other authors.

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

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At present it is generally accepted that the interaction with the two-phonon states is important for the formation of the giant resonance structure in medium and heavy spherical nuclei, which, in particular, determines their width /1-5/. However the quantitative calculations of these effects are possible within rather simple nuclear models only and were performed for a small number of nuclei. The doubly magic nuclei were calculated within the theory of finite Fermi-systems /3, 4/. In nuclei with strong pairing the calculations have been performed within the quasiparticle-phonon model /1, 5/. The latter papers have mainly investigated the electric resonances. The calculation of the M1 -resonances in the Ba , ^{140}Ce , ^{126}Te /6/ isotopes were aimed at studying the M1 -radiative strength functions in these nuclei. The influence of the two-phonon admixtures on the M1 -resonance itself has not been studied. The study of this problem for many nuclei is of great interest, especially as the simplicity of the M1 -resonance structure and its relatively low excitation energy allow one to make calculations with minimal computational simplifications. The interest to this problem is simulated also by the recent experimental data on the M1 -resonance (more likely, on the absence of it) obtained in Darmstadt /7/. In the precise experiments on the scattering of slow electrons on nuclei performed by this group, 60% of strength of the M1 -transitions predicted by the shell model has been detected in ^{58}Ni only. The M1 -resonance has not been detected neither in ^{90}Zr nor in ^{208}Pb though its existence was thought to be determined. At the same time all the calculations within the RPA (see, for instance, refs. /4, 8, 9-11/

predict the existence of the M1-resonance. A sharp weakening of the resonance M1-transitions as compared to the theoretical values obtained within the RPA is caused perhaps by two reasons. The first one is the interaction with more complex configurations primarily with the two-phonon ones, which may result in the spreading of strength of the M1-transitions over a considerable excitation energy interval. The second one, expressed in ref./7/, is a possible decrease of effective gyromagnetic factors g_s^{eff} with increasing mass number A; this will result in the predominant suppression of the M1-resonance in heavy nuclei. The influence of the continuous spectrum at the excitation energies of the M1-resonance cannot be essential.

The energy and M1-resonance excitation probabilities have been calculated for the ^{58,60}Ni, ^{90}Zr, ^{118,120}Sn, ^{124,126}Te, ^{138}Ba and ^{140}Ce nuclei. The calculations have been performed within the semimicroscopic quasiparticle-phonon nuclear model which has been thoroughly investigated in paper/12/. In the model wave function of the 1^+ state the one- and two-phonon components have been taken into account

$$\Psi_\nu(JM) = \{\sum_i R_\nu(Ji) Q^+_{JMi} + \sum_{\substack{\lambda_1\lambda_2\\ i_1 i_2}} P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)[Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2}]_{JM}\}|0>, \quad (1)$$

where $Q^+_{\lambda\mu i}$ is the phonon creation operator with the angular momentum λ, its projection μ and number i; $|0>$ is the wave function of the ground state of a doubly even nucleus (phonon vacuum). To find the energies of the states $\Psi_\nu(JM)$ and coefficients $R_\nu(Ji)$ and $P^{\lambda_2 i_2}_{\lambda_1 i_1}(J\nu)$ it is necessary to solve complicated nonlinear equations (see, for instance, ref./13/). However, as in papers/1,5,6/ we decided to calculate directly the M1-transition strength functions $b(M1,\eta)$

$$b(M1,\eta) = \frac{1}{2\pi}\sum_\nu \frac{\Delta}{(\eta-\eta_\nu)^2 + \Delta^2/4} B(M1, 0^+_{g.s.} \to 1^+_\nu). \quad (2)$$

Here

$$B(M1, 0^+_{g.s.} \to 1^+_\nu) = |\sum_i R_\nu(Ji) <0|\mathfrak{M}(M1) Q^+_{1\mu i}|0>|^2.$$

The matrix element of the magnetic dipole moment operator between the ground and i-th one-phonon 1^+-state of nucleus is calculated in the RPA. It has already been shown/1,12/ that $b(M1,\eta)$ can be calculated directly without solving complicated equations.

The Hamiltonian of the quasiparticle-phonon nuclear model includes besides the average field chosen in the form of the Saxon-Woods potential, the superconducting pairing interaction and the so-called long-range forces which generate in the nucleus the phonon excitations. These forces are chosen in the factorized form, which allows one to simplify the calculations essentially. The phonons with $I,\pi = \lambda, (-1)^\lambda$ are generated by the multipole forces whereas the phonons with $I,\pi = \lambda, (-1)^{\lambda+1}$ by the spin-multipole forces. The parameters of the single-particle potential chosen on the basis of the results of refs./14/, are given in *Table 1*. The single-particle energies, wave functions and matrix elements of the multipole and spin-multipole operators are calculated by using the code REDMEL which realizes the numerical method of solving the spherically symmetrical Schrodinger equation suggested in paper/15/.

Table 1

Parameters of the Saxon-Woods Potential and Pairing Interaction Constants

A	N,Z	r_0, fm	V_0 MeV	$\varkappa$ fm^2	α fm^{-1}	$G_{N,Z}$ MeV
59	N=31	1.31	46.2	0.413	1.613	0.280
	Z=27	1.24	53.7	0.308	1.587	0.302
91	N=53	1.29	44.7	0.413	1.613	0.168
	Z=39	1.24	56.9	0.338	1.587	0.194
121	N=71	1.28	43.2	0.413	1.613	0.122
	Z=51	1.24	59.9	0.346	1.587	0.136
127	N=73	1.28	43.4	0.413	1.613	0.124
	Z=53	1.24	59.7	0.350	1.587	0.130
141	N=83	1.27	46.0	0.413	1.613	0.116
	Z=59	1.24	57.7	0.349	1.587	0.122

As a result the single-particle spectra used differ from those obtained earlier in paper [16] by using the semi-analytic way of solving the Schrödinger equation[17]. A new spectrum has a larger density of the single-particle levels and the wave functions of the quasi-stationary states have a correct asymptotic form what is more important. The constants of the pairing forces G_N, G_Z have been determined by a usual procedure from the experimental values of the pairing energies. The values of $G_{N,Z}$ are also given in *Table 1*. It is more difficult to choose the constants of the long-range forces, the isoscalar and isovector ones. For the dipole forces they have been determined from the position of the giant dipole resonance in each nucleus and from the condition $\omega_{11}=0$, this allows one to exclude with good accuracy the influence of a spurious state which is caused by the breaking of the translational invariance[19]. The constants of the quadrupole and octupole interactions are chosen so as to fit the theoretical values of the 2_1^+ and 3_1^- level energies to the experimental ones. The theoretical energies have been calculated taking into account the admixture of the two-phonon components to their structure[16]. The values of the multipole constants $\lambda > 3$ are chosen to be less by an order of 1.5 than in monograph[20]. At these values of the constants the low-lying collective states with $\lambda > 3$ do not appear. The isovector constants $\kappa_1^{(\lambda)}$ have been calculated by the isoscalar $\kappa_0^{(\lambda)}$ ones using the relation:

$$\kappa_1^{(\lambda)} = -\kappa_0^{(\lambda)} \times 0{,}2(2\lambda + 3)\,. \qquad (3)$$

The relation (3) is a renormalized version of formula (6.127) from the monograph[20]. It provides such a value of the isovector quadrupole constants $\kappa_1^{(2)}$ that the position of the E2-resonance with $T=1$ coincides with the experimental one ; this value is close also to the estimates of the hydrodynamic model. The value of the constants of the spin-spin interaction has been studied in paper [8]. It has been obtained that the values

$$\kappa_0 = 0, \quad \kappa_1 = -4\pi \times 28/A \ \text{MeV} \qquad (4)$$

provide the position of the M1-resonance in the RPA, which is in agreement with a scarce experimental data. This value of the constants agrees with the data of the Copenhagen group[21]. Alongside with (4) in some nuclei the calculations have been performed for other values of κ_0, κ_1. For the spin-multipole forces with $\lambda > 1$, we assume

$$\kappa_0^{(\lambda)} = 0, \quad \kappa_1^{(\lambda)} = -\frac{4\pi \times 28}{A < r^{\lambda-1} >^2} \ \text{MeV} \times \text{fm}^{-2\lambda+2}\,.$$

The value of the effective gyromagnetic factors were taken from paper [8]: $g_s^{eff} = 0.8\, g_s^{free}$, $g_\ell^n = 0$, $g_\ell^p = 1.0$. For the parameter Δ we took the value of 0.1 *MeV*. Such a value of Δ allows one to distinguish between individual states with large value of B(M1) and does not result in additional broadenning of the M1-resonance.

Since the one-phonon excitations can be collective and two-quasiparticle, the two-phonon part of the wave function (1) may include the components violating the Pauli principle. To weaken this effect an additional requirement was imposed on the two-phonon part of the wave function (1), that is one of the phonons $Q^+_{\lambda_1\mu_1 i_1}$ or $Q^+_{\lambda_2\mu_2 i_2}$ should necessarily be collective. Of course, this method is not consistent, but a consistent consideration of the Pauli principle complicates strongly the problem.

According to the influence of the two-phonon admixtures on the M1-resonance the nuclei studied are divided into two groups: The first group includes the nuclei with magic number of neutrons in which the influence of the quasiparticle-phonon interaction on the M1-resonance is weak; the second group includes all the rest nuclei (Ni , Sn , Te) where it influences strongly the distribution of strength of the M1- transitions in the resonance region. The first group includes ^{90}Zr, the resonance of which has been studied by many authors[4,8,10]. We shall begin to discuss our results with this nucleus. *Figure 1* shows the results of our calculations of the M1-resonance in the RPA *(a)-c))* and of other authors[4,10] *(d)-f))*. Our calculations are given for three values of the constant κ_0: a) $\kappa_0 = 0$; b) $\kappa_0 = 0.5\kappa_1$; c) $\kappa_0 = 0.9\kappa_1$. With increasing $|\kappa_0|$ the states with larger values of B(M1) become close to each other; this is mainly due to the increasing excitation energy of the lowest state. The calculations shown in *Fig. 1d)-f)* are performed within the theory of finite Fermi-systems. *Figure 1d)* is taken from paper [4] which did not consider the pairing interaction in the proton system. Therefore, the M1-resonance comprises only one neutron particle-hole state $1g_{7/2}-1g_{9/2}$. The other two *figures (1e)* and *1f))* show

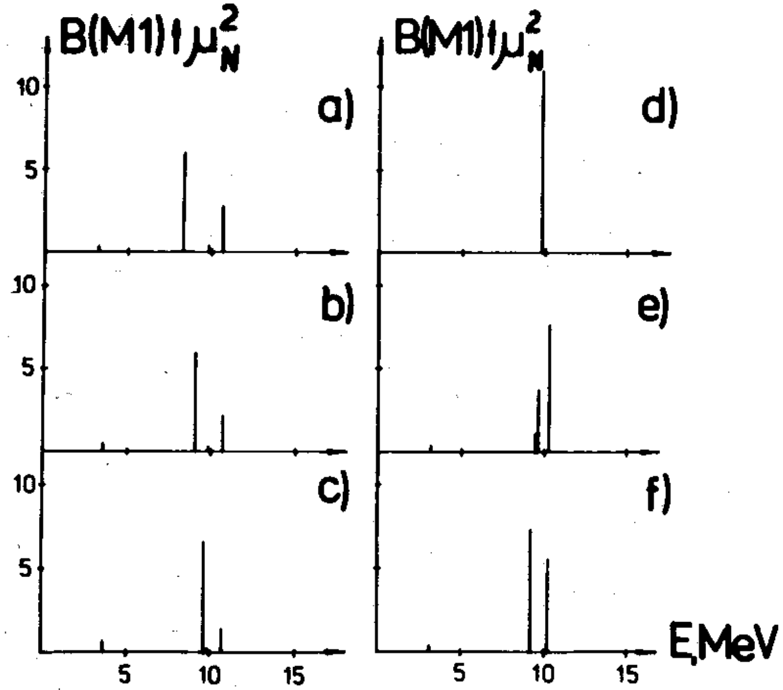

Fig. 1. The M1*-resonance in* ^{90}Zr *in the RPA. Our calculations with the effective spin-spin-forces* $\kappa_1 = -28\times 4\pi/A$ MeV *a)* $\kappa_0 = 0$ *; b)* $\kappa_0 = 0.5\kappa_1$ *; c)* $\kappa_0 = 0.9\kappa_1$ *; Calculations by the theory of finite Fermi-systems: d) ref.* [4] *; e)-f) ref.* [9].

the results of paper [10] for two values of factor $(dn/d\epsilon_F)^{-1}$ entering into the expression for the effective particle-hole interaction. We give these results in order to show that different semimicroscopic models in the RPA give similar results for the M1-resonance at a certain choice

of the strength parameters. The total probability of the M1 -resonance excitation differs in various papers mainly due to different values of g_s^{eff}. The interaction with the two-phonon states influences the M1-resonance in ^{90}Zr slightly. As it is seen from *Fig. 2a)* which shows the function $b(M1, \eta)$ in ^{90}Zr for $\kappa_0 = 0.9\kappa_1$, the region of location of the M1-resonance is the same with the mean energy being somewhat lower. In the calculations by the theory of finite Fermi-systems [4] (*fig. 2b*) the two-phonon states influence the M1-resonance stronger. The only 1^+ state obtained in the RPA (*fig. 1d*) spread into three states which are located in the interval $\Delta E \approx 2.5$ *MeV*; besides there appeared a weak state with $E = 15.5$ *MeV*. However we should like to mention that the "final" form of the M1-resonance in paper [4] and in our calculations is almost the same. Thus, neither our results nor the

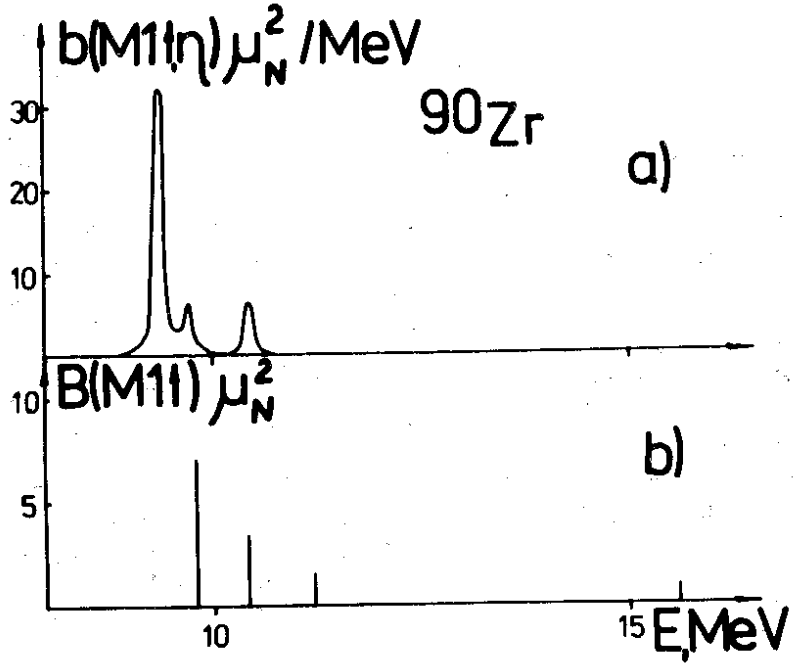

Fig. 2. The M1*-resonance in* ^{90}Zr *taking into account the interaction with the two-phonon states: a) our calculations; b) calculations of paper* [4].

results of paper /4/ cannot explain by the interaction with the two-phonon configurations the "disappearance" of a large portion of strength of the M1-transitions in ^{90}Zr which was observed in ref. /7/. In other nuclei with magic number of neutrons the influence of the two-phonon admixtures on the structure of the M1-resonance is negligible too.

Rather reliable experimental data on the M1-resonance exist for the ^{58}Ni nucleus /22/. In the experiments on the inelastic scattering of electrons in this nucleus there were observed two groups of states with the excitation energies from 6.0 to 8.5 *MeV* ($\Sigma B(M1)\uparrow \simeq 2.7\mu_N^2$) and from 9.8 to 11 *MeV* ($\Sigma B(M1)\uparrow \approx 5.7\mu_N^2$). It if confirmed that they are the levels with $I^\pi = 1^+$, though the group with lower excitation energy perhaps comprises the 1^+ and 2^- levels. The calculations of the M1-resonance in ^{58}Ni with the isovector spin-spin forces in the RPA (*Fig. 3a*), (see also ref. /8/) predict the existence of one 1^+ level with $\omega = 10$ *MeV* and $B(M1)\uparrow \simeq 10\mu_N^2$ (at $g_s^{eff} = 0.8g_s^{free}$). The inclusion of the isoscalar spin-spin forces leads to the splitting of this state into two states spaced by the distance of about 200 *KeV (fig. 3c)* and with the same total value of $B(M1)\uparrow$. The interaction with the two-phonon states influences strongly the M1-resonance in ^{58}Ni; this results in the spreading of strength of the only one-phonon state with a large value of $B(M1)$ in the interval $\Delta E \approx 2$ *MeV*. The maximum of the function $b(M1,\eta)$ is by 400 *keV* lower than the one-phonon 1^+ state itself (*Fig. 3a*). The position of the maximum $b(M1,\eta)$ is almost independent of $|\kappa_0|$, as it can be seen from *Fig. 3b), c)*, but the distribution of strength of the M1-transitions in the resonance region changes noticeably. At the values of constants (4) the maximum of $b(M1,\eta)$ is by 600 *keV* lower than the experimentally observed group of 1^+-states with strong M1-transitions /22/. The total probability of excitation $\Sigma B(M1)\uparrow$ in the interval $\Delta E = 1$ *MeV* which is symmetric with respect to the maximum $b(M1,\eta)$ is $7.8\mu_N^2$. By renormalizing the constant κ_1 we can make the maximum $b(M1,\eta)$ coincide with the experimental energy. *Figure 4a)-c)* shows how the form of the M1-resonance changes with increasing $|\kappa_1|(\kappa_0=0)$.

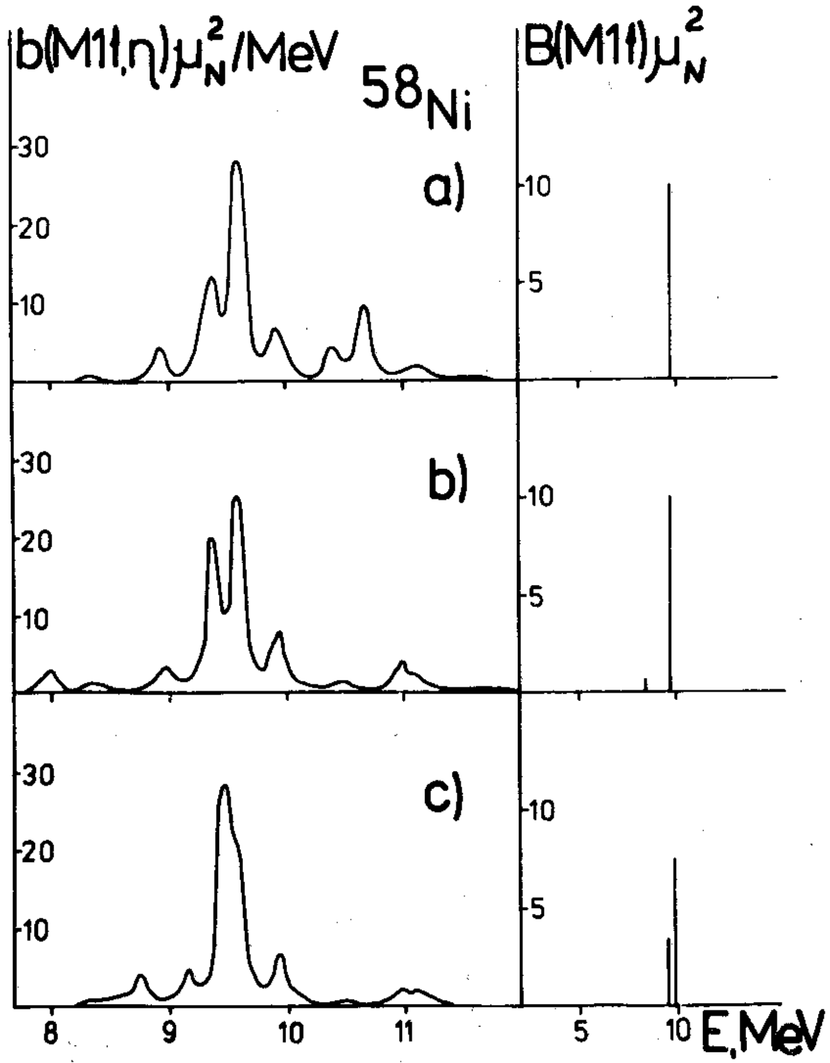

Fig. 3. The M1-*resonance in* ^{58}Ni. *The right part is the calculation in the RPA. The left part if the strength function* $b(M1,\eta)$ *taking into account the interaction with the two-phonon states.* $\kappa_1 = -28\times 4\pi/A$ MeV a) $\kappa_0 = 0$, *b)* $\kappa_0 = 0.5\kappa_1$, *c)* $\kappa_0 = 0.9\kappa_1$.

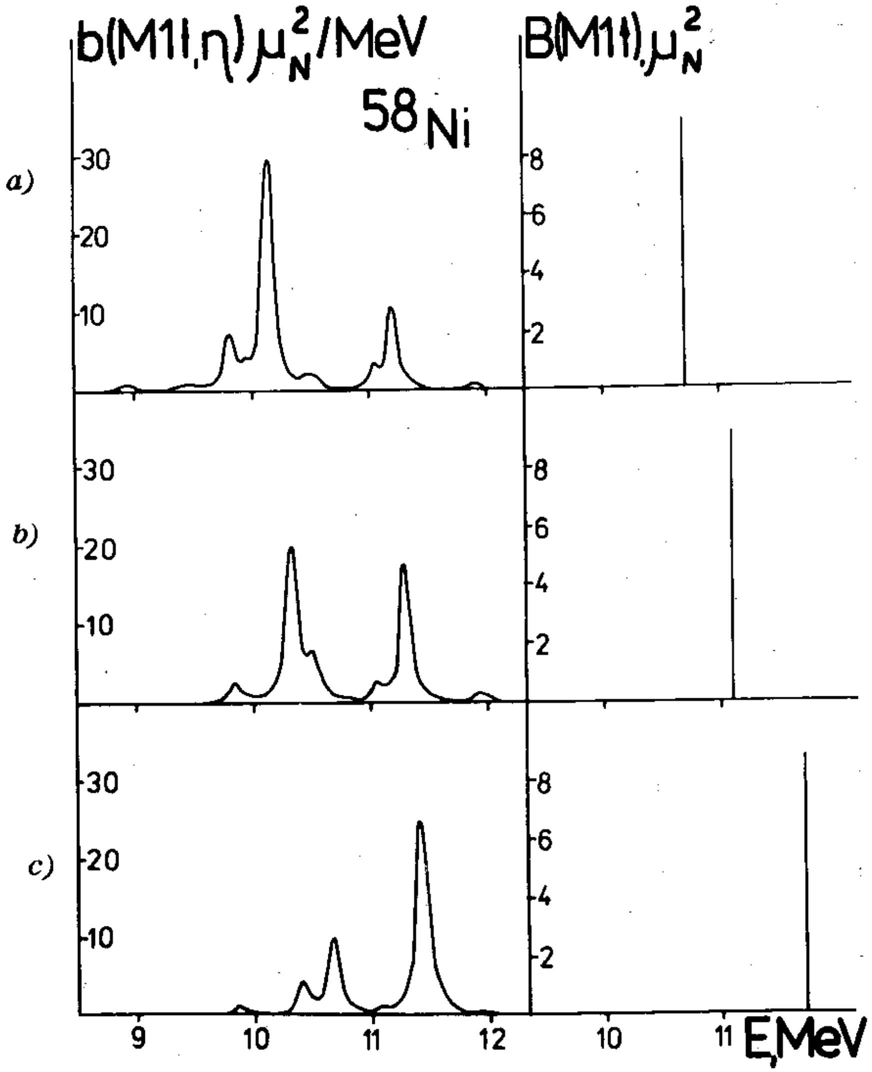

Fig. 4. M1-*resonance in* ^{58}Ni. *The right part is the calculation in the RPA. The left part is the strength function* b(M1, η) *taking into account the interaction with the two-phonon states.* $\kappa_0 = 0$ *a)* $\kappa_1 = -37\times 4\pi/A$ MeV, *b)* $\kappa_1 = -42\times 4\pi/A$ MeV, *c)* $\kappa_1 = -50\times 4\pi/A$ MeV.

The best agreement with the experimental excitation energy is achieved at $\kappa_1 = -37\times 4\pi/A$ MeV. As it is seen from *Fig. 4* the increase in $|\kappa_1|$ results in increasing energy of the main peak of the M1-resonance, but its amplitude decreases. As a result at the value $\kappa_1 = -37\times 4\pi/A$ MeV in the interval $\Delta E = 1$ *MeV* with the center at the maximum b(M1, η) the sum value of B(M1) is $5.9\mu_N^2$ which is in good agreement with the experimental values. However, we fail to obtain the 1^+ states in the region of 6-8 *MeV* with the value of B(M1) close to the experimental one at any values of κ_0 and κ_1. The only one-phonon 1^+ state in this region has B(M1) $\simeq 0.5\mu_N^2$, and the interaction with the two-phonon states does not change this value. For the states in this region B(M1) increases in heavier isotopes of Ni only; this is due to the strengthening of the M1-transition from the two-quasiparticle neutron 1^+ state $2p_{3/2}-2p_{1/2}$ in these nuclei*. This result testifies to the assumption that the lower group of states in ^{58}Ni comprises mainly the states with $I^\pi = 2^-$. To make more profound conclusions it is necessary to study the spectrum of the 2^- states in this region of excitation energies.

It is seen from *Figs. 3a)-c)* that the form of the M1-resonance in ^{58}Ni changes depending on the ratio of the constants κ_0 and κ_1. A more interesting situation occurs when in the one-phonon approximation the M1-resonance is formed by several states with slightly different values of B(M1). As an example *Figs. 5a)-c)* and *Fig. 6* show the M1 -resonance in the ^{118}Sn and ^{124}Te nuclei. In ^{118}Sn the M1-resonance is given for three values of κ_0: a) $\kappa_0 = 0$, b) $\kappa_0 = 0.5\kappa_1$, c) $\kappa_0 = 0.9\kappa_1$ With increasing $|\kappa_0|$ the two strongest one-phonon 1^+ states become close to each other and at the same time B(M1) decreases for the states with higher excitation energy (the right part of *Fig. 5a)-c)*. The interaction with the two-phonon states leads to that two peaks approach each other very slowly, and the integral values of B(M1) in them become close to each other. There occurs redistribution of strength of the M1-transitions in

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* The state $(2p_{1/2}, 2p_{3/2})^N_{1^+}$ in the Ni isotopes is at an energy of about 6.5 *MeV*, but beginning from the ^{60}Ni isotope it becomes the particle-hole state; this is just the reason of strengthening of B(M1).

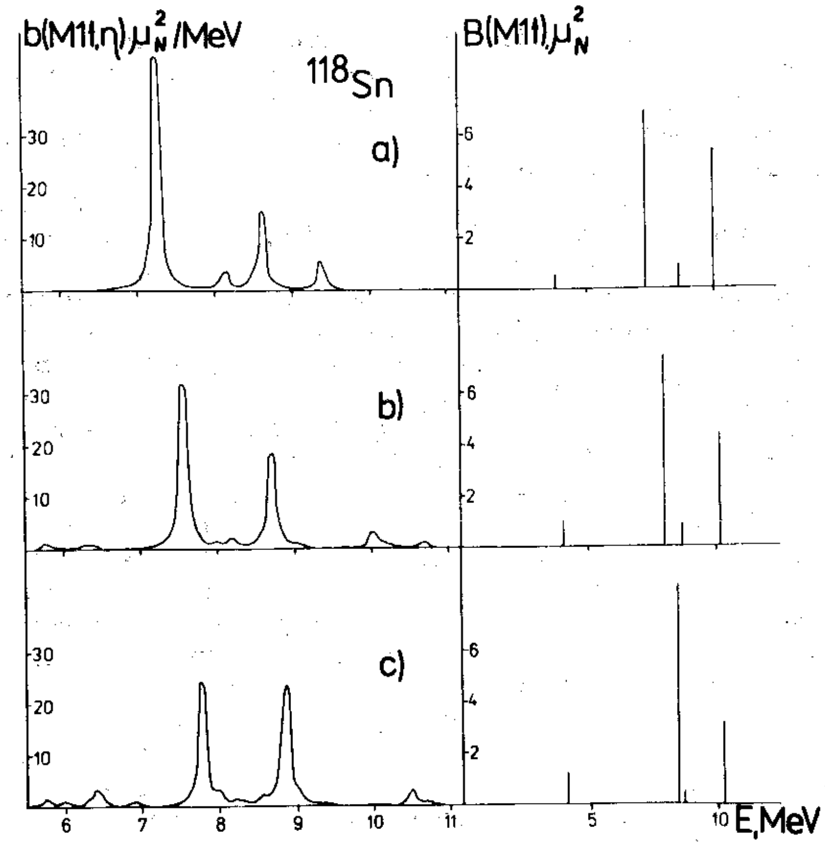

Fig. 5. M1-*resonance in* ^{118}Sn. *The right part is the calculation in the RPA. The left part is the strength function* b(M1, η) *taking into account the interaction with the two-phonon states* $\kappa_1 = -28{\times}4\pi/A{\cdot}MeV$ a) $\kappa_0 = 0$, b) $\kappa_0 = 0.5\kappa_1$, c) $\kappa_0 = 0.9\kappa_1$.

the region of about 1.5 *MeV* occupied by the M1-resonance. In ^{124}Te (*Fig. 6*) the M1-resonance in the RPA are represented by three 1^+ states with almost equal B(M1). The two-phonon admixtures influence mainly the states with larger excitation energy, which are fragmented strongly. The corresponding strength of the M1-transitions is spread in the interval of more than 2 *MeV*. The state at an energy of 7 *MeV* almost does not change neither its position nor the value of B(M1). This is perhaps due to a larger density of the two-phonon 1^+ states at an energy of about 10 *MeV*.

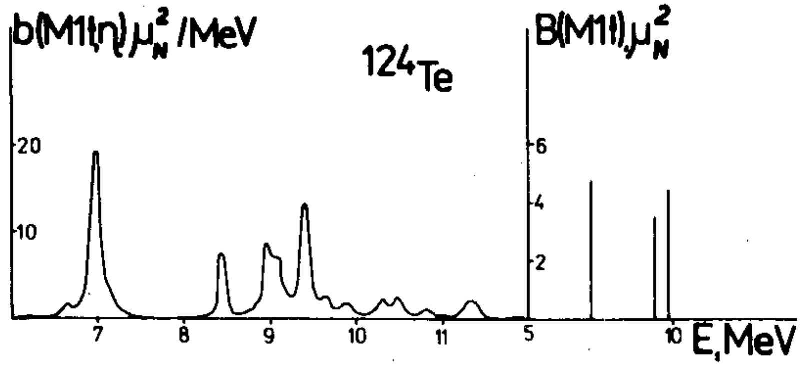

Fig. 6. M1-*resonance in* ^{124}Te. *The right part is the calculation in the RPA. The left part is the strength function* b(M1, η) *taking into account the interaction with the two-phonon states* $\kappa_0 = 0$, $\kappa_1 = -28\times 4\pi/A$ MeV .

One more source of information about the distribution of strength of the M1-transitions in nuclear spectra is the radiative strength function k(M1). At the energies close to the neutron binding energy B_n the value of k(M1) is obtained from the data on the (n,γ) and (γ, n) reactions. At other energies it is determined from the inelastic scattering of γ-quanta on nuclei. It is true that in the (γ, γ') reaction the values of k(M1) are overestimated in comparison with the (n, γ) and (γ, n) data; this is caused by the mechanism of inelastic scattering of γ-quanta in nuclei, which separates the 1^+ states with the largest values of B(M1) in a narrow energy region. Systematics of the experimental data on k(M1) in different nuclei [23] gives an average value of $\langle k(M1)\rangle \simeq (15\text{-}20)\times 10^{-9}$ MeV^{-3}; however, it exceeds greatly the single-particle values. It was assumed that a sharp increase in k(M1) in nuclei with $A \simeq 140$ [24, 25] observed experimentally is due to the proximity of the M1-resonance to B_n in these nuclei. This assumption was verified by the theoretical calculations within the quasiparticle-phonon nuclear model [6]. It should be noted that a consistent microscopic calculation of k(M1) in spherical nuclei is possible if the fragmentation of the one-phonon states over the two-phonon ones is taken into account because of the small density of one-phonon 1^+ levels and large differences in value of B(M1) among them.

With the known strength function of the M1-transitions $b(M1, \eta)$, the value of $\langle k(M1)\rangle$ can easily be calculated by the following formulae:

$$\langle k(M1)\rangle = \sum_{\nu \in \Delta E} \Gamma_{\gamma_0}(M1, \eta_\nu)/E_\gamma^3 \Delta E$$

$$\sum_{\nu \in \Delta E} \Gamma_{\gamma_0}(M1, \eta_\nu) = 3.76 \times 10^{-3} \int_{E-\Delta E/2}^{E+\Delta E/2} \eta^3 b(M1, \eta)\, d\eta \ (eV),$$

where Γ_{γ_0} is in *eV*, E_γ and ΔE in *MeV*. The calculated values for most of the nuclei studied are given in *Table 2*. The calculations have been performed with κ_1 calculated by (4) for three values of κ_0 : $\kappa_0 = 0$; $\kappa_0 = 0.5\kappa_1$; $\kappa_0 = 0.9\kappa_1$. The averaging interval $\Delta E = 1$ *MeV* (the discussion of the dependence of results on ΔE is given in ref. /6/). For those nuclei which have no experimental data (a majority of them), we give the values of $\langle k(M1)\rangle$ at B_n . We should like to point the difference of our results and those of ref. /6/ for ^{140}Ce and ^{138}Ba. This difference is caused by a slight difference in the single-particle schemes and superfluid parameters. In ^{140}Ce the M1-resonance is slightly fragmented but at the same time it is very close to B_n . Thus a slight change of its position towards B_n results in a considerable change of $\langle k(M1)\rangle$. So, at $\kappa_0 = 0$ the center of mass of the M1-resonance in this nucleus is at an energy of 8.4 *MeV* (the experimental value is 8.7 *MeV* /26/), and therefore the value of $\langle k(M1)\rangle$ is too low in contrast with experiment. The increasing $|\kappa_0|$ leads to the increase in the resonance energy and in $\langle k(M1)\rangle$. The same is in ^{140}Ce with increasing $|\kappa_1|$. The experimental and theoretical values of the resonance energy coincide at $\kappa_1 = -34 \times 4\pi/A$ MeV and $\langle k(M1)\rangle = 25 \cdot 10^{-9} MeV^{-3}$. It should be noted that the value of B(M1) for the M1-resonance in ^{140}Ce is smaller in our paper than the experimental one (the latter is known with an accuracy of about 50% /26/); this causes

Table 2

M1-Radiative Strength Functions

Nuclei	E_γ MeV	$\langle k(M1)\rangle \cdot 10^9$ MeV^{-3} Exper.	Theory $\varkappa_0 = 0$	Theory $\varkappa_0 = 0.5\varkappa_1$	Theory $\varkappa_0 = 0.9\varkappa_1$
^{118}Sn	9.33	-	4.5	2.6	9.5
^{120}Sn	9.11	-	8.9	9.5	17.8
^{124}Te	9.41	-	18.2	14.9	11.3
^{126}Te	7.915	39	3.0	18.0	24.0
	9.09	-	30	22	15
^{138}Ba	8.54	90±35	44	37	5.6
^{140}Ce	9.1	37	7.7	35	31

the low value of $\langle k(M1)\rangle$. A high experimental value of k(M1) at the energy 7.915 *MeV* in ^{126}Te is obviously due to the fact that it is obtained in the (γ, γ') reaction /27/. By increasing $|\kappa_1|$ in this nucleus, we cannot essentially change the theoretical value of $\langle k(M1)\rangle$ which is determined mainly by the distribution of strength of the M1-transitions in the resonance region with width of about 2 *MeV*. Though the energies of the M1-resonance and the values of B_n in the Sn and Te isotopes are almost the same, the value of the radiative strength function is less in the Sn isotopes. The reason of this is a weak fragmentation of the M1-resonance in the tin isotopes. Since in the ^{58,60}Ni and ^{90}Zr nuclei the M1-resonance is by 2-3 *MeV* lower than B_n, the values of $\langle k(M1)\rangle$ at B_n are small and we did not give them in *Table 2*.

The results of this paper demonstrate that in many nuclei the two-phonon admixtures influence greatly the distribution of the M1-transitions in the region of the magnetic dipole resonance although this influence is not strong, and we cannot speak about the "disappearance" of the resonance. Though very simple the separable spin-spin interaction at a certain choice of constants allows one to describe fairly well the experimental data. A weak point of these calculations and of all model calculations with effective forces is the uncertainty of the parameters (strength constants, effective charges and so on). So, for the constants of the isovector spin-spin forces one should obviously use the value $\kappa = -35\times 4\pi/A$ MeV in contrast with that of paper /8/ (see formula (4)) obtained in the RPA. In fact uncertain is the choice of κ_0. The values of the effective gyromagnetic factors allow one to describe B(M1) well in ^{58}Ni, but in ^{140}Ce we have obtained too low values for B(M1) (the same in other papers /9,10/). A very large experimental value of B(M1) for the M1-resonance in ^{140}Ce (in theoretical calculations it is reproduced under the assumption that $g_s^{eff} = g_s^{free}$ only) contradicts the assumption of paper /7/ about the reduction of effective g_s factors with increasing A. The calculation of the radiative strength functions with the same parameters as for the M1-resonances, gives a satisfactory agreement with experiment. It should be pointed once more that for these quantities the consideration of the influence of the two-phonon states is of great importance.

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REFERENCES

1. *Soloviev V.G., Stoyanov Ch., Vdovin A.I. Nucl. Phys., 1977, A288, p.376.*
2. *Румянцев Б.А. ЯФ, 1976, 24, с.1098.*
   Faessler A. Selected Topics in Nucl. Struct., 1976, JINR, D-9920, v.2, p.242, Dubna, 1976.
   Камерджиев С.П. Изв. АН СССР, сер. физ., 1977, 41, с.1230.
3. *Dehesa J.S., Speth J., Faessler A. Phys. Rev.Lett., 1977, 38, p.208.*
4. *Dehesa J.S. Ph.D.Thesis. Rheinischen Friedrich-Wilhelms Universitat zu Bonn, 1977.*
5. *Soloviev V.G., Stoyanov Ch., Voronov V.V. Nucl. Phys., 1978, A304, p.503.*
6. *Воронов В.В., Соловьев В.Г., Стоянов Ч. ОИЯИ, Р4-11591, Дубна, 1978.*
   Ponomarev V.Ju., et al. Contributed papers for III International Symposium on Neutron Capture Gamma-Ray Spectroscopy and Related Topics, BNL, No. 63, 1978.
7. *Knüpfer W. et al. Phys. Lett., 1977, 77B, p.367.*
8. *Вдовин А.И., Стоянов Ч., Юдин И.П. ОИЯИ, Р4-11081, Дубна, 1977.*
9. *Борзов И.Н., Камерджиев С.П. Изв. АН СССР, сер. физ., 1977, 41, с.4. Ткачев В.Н., Борзов И.Н. Камерджиев С.П. ЯФ, 1976, 24, с.715. Борзов И.Н., Ткачев В.Н. Изв. АН СССР, сер. физ., 1977, 41, с.1263.*
10. *Борзов И.Н., Камерджиев С.П. Препринт ФЭИ-580, Обнинск, 1975.*
11. *Cwick S., Wygonowska M. Acta Phys. Pol., 1973, B4, p.223.*
12. *Соловьев В.Г. Теория сложных ядер. "Наука", М., 1971. Соловьев В.Г. ЭЧАЯ, 1978, 9, с.810.*
13. *Вдовин А.И., Кырчев Г., Стоянов Ч. ТМФ, 1974, 21, с.137.*
14. *Чепурнов В.А. ЯФ, 1967, 6, с.955.*
15. *Bang J. et al. Nucl.Phys., 1976, A261, p.59.*
    Гиззаткулов М.Х., Пузынин И.В., Ямалеев Р.М. ОИЯИ, Р11-10029, Дубна, 1976.
16. *Вдовин А.И., Стоянов Ч. Изв. АН СССР, сер. физ., 1974, 38, с.2598, 2604.*
17. *Ширикова Н.Ю. ОИЯИ, Р4-3712, Дубна, 1968.*
18. *Вдовин А.И., Комов А.Л., Малов Л.А. ОИЯИ, Р4-5125, Дубна, 1970.*
19. *Petersen D.F., Veje C.J. Phys. Lett., 1967, 24B, p.449. Pyatov N.I., Salamov D.I. Nucleonica, 1977, 22, p.127.*
20. *Бор О., Моттельсон Б. Структура ядра. "Мир", М., 1977, т.2.*
21. *Castel B., Hamamoto I. Phys. Lett., 1976, 65B, p.27.*
22. *Lindgren R.A. et al. Phys. Rev., 1976, C14, p.1789.*
23. *Bollinger L.M. Proc. of the Int. Conf. on Photonucl. React. and Appl., Asilomar, 1973, ed. by B.Berman (Lawrence Livermore Lab., Univ. of California, 1973), p.783.*
24. *Chrien R.E. et al. Phys. Rev., 1974, C9, p.1622. Holt R.J., Jackson H.E. Phys. Rev., 1975, C12, p.56.*
25. *Laszewski R.M., Holt R.J., Jackson H.E. Phys. Rev., 1976, C13, p.2257.*
26. *Pitthan R., Walcher Th. Zeitsch. Naturforsch., 1972, 27a, p.1683.*
27. *Wolf A., Moreh R., Shahal O. Nucl. Phys., 1974, A227, p.373.*

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