объединенный
институт
ядерных
исследований
дубна

5865/2-80

8/12-80

E4-80-546

V.Navrotska-Rybarska, O.Stoyanova,
Ch.Stoyanov

INFLUENCE OF GROUND STATE CORRELATIONS ON THE PROPERTIES OF THE FIRST 2^{+} AND 3^{-} STATES IN SOME Sm ISOTOPES

Submitted to ЯФ

1980

1. INTRODUCTION

The models describing excited nuclear states at intermediate and high excitation energies have been developed in recent years [1-3]. In most of these models the RPA is used as the basic approximation. There arises a guestion whether one and the same set of parameters can be used for the description of the properties of low-lying and high-lying nuclear excitations. It can definitely be answered upon investigating more thoroughly the accuracy of the RPA and the values of corrections to this approximation, which should be taken into account in describing the low-lying excited nuclear states.

The accuracy of the RPA is being investigated more than ten years. The terms of the nuclear Hamiltonian causing nonlinear kinematic and dynamical effects were investigated [4]. One of the recent papers [4] gives a review of the most important papers devoted to this problem.

In recent years the quasiparticle-phonon model [1] is of great interest. This model allowed one to calculate the spreading widths of the giant E1 and E2 resonances [5], the influence of these resonances on the radiative strength functions [6], enhancement of M1 transitions in some spherical nuclei [7], and the regions of location of M2 transitions [8]. Some interesting results have also been obtained for the deformed nuclei [9].

An essential advantage of the quasiparticle-phonon model is a large phonon space. According to the above-mentioned papers, to adequatly describe the distribution of few-quasiparticle components of the nuclear wave function, one should take into account numerous collective and noncollective phonons with momentum from 1 to 7 and more of negative and positive parity. The structure of these phonons is calculated within the RPA. In refs. [10,11] the accuracy of this approximation in spherical and transitional nuclei is investigated for the case when a large phonon space is used. It is studied [10] that the accuracy of the RPA in spherical nuclei is good for the phonons of any multipolarity in nuclei with one closed shell and in the adjucent nuclei. However, with increasing collectivity of the first 2^+ and 3^- states, the number of quasiparticles in the ground state of an even-even nucleus increases and the accuracy of the RPA becomes worse. The simplest way of taking into account the influence of quasiparticles in the ground state on the properties of excited states has been suggested in refs.[12,13]. The correlations in the ground states have been considered in more detail in ref.[14]. The coupling between quasiparticle and phonon branches of excitation arising due to the correlations in the ground state was also considered in that paper. Using the variational principle a system of equations connecting the quasiparticle and phonon degrees of freedom was obtained. However, these effects were not estimated numerically.

In this paper, based on the results of ref.[14], we estimate the influence of correlations in the ground state on the energies and probabilities of electric transitions of the first excited 2^+ and 3^- states. The calculations have been performed using a large basis of single-particle states; this makes it possible to reduce to minimum the influence of the effective charge on the final result. The renormalization of parameters of multipole forces in order to make the agreement with the experimental data better is discussed.

II. STATEMENT OF THE PROBLEM AND BASIC EQUATIONS

Let us proceed from the nuclear model whose Hamiltonian includes the average field $\mathcal{H}_0$, pairing $\mathcal{H}_P$ and multipole residual forces $\mathcal{H}_{QQ}$. Sometimes the Hamiltonian includes also the spin-multipole forces[1]. We use the Hamiltonian which does not include these forces; however their influence on the final result will be discussed below. Now we write the Hamiltonian as follows:

$$\mathcal{H} = \mathcal{H}_0 + \mathcal{H}_P + \mathcal{H}_{QQ}, \tag{1}$$

where

$$\mathcal{H}_0 = \sum_{jm} (E_j - \lambda) a^+_{jm} a_{jm}, \tag{2}$$

$$\mathcal{H}_P = -\frac{G}{4} P^+ P,$$

(3)

$$P^+ = \sum_{jm} (-1)^{j-m} \; a^+_{jm} \, a_{j-m}, \tag{4}$$

$$\mathcal{H}_{QQ} = -\frac{1}{2} \sum_{LM} \chi_L \hat{Q}^+_{LM} \hat{Q}_{LM}, \tag{5}$$

$$\hat{Q}_{LM} = \sum_{\substack{j_1 m_1 \\ j_2 m_2}} \langle 1 | Q_{LM} | 2 \rangle \, a^+_{j_1 m_1} a_{j_2 m_2}, \tag{6}$$

$$Q_{LM} = r^L Y_{LM} (\theta, \phi), \tag{7}$$

$$\langle 1 | Q_{LM} | 2 \rangle = \int \psi^+_1(x) \, Q_{LM}(x) \, \psi_2 (x) \, dx , \tag{8}$$

and $\psi_1 (x)$ are the single-particle wave functions.

Using the Wigner-Eckart theorem, the matrix element (8) can be written as follows:

$$\langle 1 | Q_{LM} | 2 \rangle = \frac{\langle 1 || Q_L || 2 \rangle}{\sqrt{2\Omega_1}} \langle 2, LM | 1 \rangle ,$$

where $\langle 2, LM | 1 \rangle$ is the Clebsch-Gordan coefficient.

Now we transform the Hamiltonian (1) passing to quasi-particles

$$\begin{aligned} \alpha^+_{jm} &= u_j a^+_{jm} + (-1)^{j-m} v_j a_{j-m} , \\ \alpha_{jm} &= u_j a_{jm} + (-1)^{j-m} v_j a^+_{j-m} \end{aligned} \tag{9}$$

under the condition that

$$u_j^2 + v_j^2 = 1 .$$

After the transformation (9) the Hamiltonian is a quadratic form of the following operators:

$$A^+_j, A_j, B_j, A_{LM}[jj'], \quad A^+_{LM}[jj'], \quad B_{LM}[jj'],$$

where

$$A^+_j = \frac{1}{\sqrt{4\Omega}} \sum_m (-1)^{j-m} \alpha^+_{jm} \alpha^+_{j-m} = (A_j)^+ ,$$

$$B_j = B^+_j = \frac{1}{\sqrt{2\Omega_j}} \sum_m \alpha^+_{jm} \alpha_{jm} ,$$

$$2\Omega_j = 2j + 1 ,$$

$$A^+_{LM}[12] = \frac{1}{\sqrt{2}} \sum_{m_1 m_2} <1, 2|LM> a^+_1 a^+_2 = (A_{LM}[12])^+ , \tag{10}$$

$$B^+_{LM}[12] = B_{LM}[12] = \sum_{m_1 m_2} <1,-2|LM> a^+_1 a_2 .$$

The commutation relation of the operators (10) has a complex form, for instance

$$[A_{LM}[22'], A^+_{\lambda\mu}[11']] = \delta_{\lambda L}\,\delta_{\mu M}\, L(B^+_{L'M'}[33']) , \tag{11}$$

where $L(P)$ denotes the linear form with respect to the operator P . Let us denote the phonon creation operator as follows:

$$C^+_n = \sum_1 (a_{n1} A^+_1 - b_{n1} A_1) ,$$

$$C^+_{L\mu}[n] = \sum_{11'} (\psi^n_L [11'] A^+_{L\mu}[11'] - \phi^n_L [11'] (-1)^{L-\mu} A_{L-\mu}[11']) , \tag{12}$$

Let $|\,>$ be a phonon vacuum. It is assumed in the RPA that the right-hand side of the commutator (11) averaged over the phonon vacuum can be substituted by $\delta_{\lambda L}\delta_{\mu M}$. This means that the following relation

$$<|B^+_{LM}[11]|> = 0 ,$$

$$<|B^+_j|> = 0 \tag{13}$$

should hold.

Thus, the operators A^+ and A as C^+ and C will satisfy the boson commutation relations [15].

As has been shown in ref.[10] the assumption (13) is fulfilled for a small number of nuclei only. Following papers [12,13] we shall take into account in very simple way the right-hand side of the commutation relations (11).

Let us introduce the following quantities:

$$\rho_j = \frac{1}{\sqrt{2\Omega_j}} \langle | B_j^+ | \rangle,$$

$$\frac{1}{\sqrt{2\Omega_j}} \langle | B_{LM}^+[11'] | \rangle = \delta_{L0}\delta_{11'}\rho_1 \quad . \tag{14}$$

Thus, relation (13) will not be fulfilled. Let also the operators C_n^+, C_n and $C_{LM}^+[n]$, $C_{LM}[n]$ satisfy the boson commutation relations. Express the Hamiltonian through C_n and $C_{LM}[n]$ and construct the functional

$$\mathcal{L} = \langle | H | \rangle - \sum_1 \mu_1 (v_1^2 + u_1^2) -$$

$$- \sum_{nL11'} \omega_L^n (1-\rho_{11'})(\psi_L^n[11']\psi_L^n[11'] - \phi_L^n[11']\phi_L^n[11']) - \tag{15}$$

$$- \sum_{1n} \omega_0^n (1-\rho_1)(a_{n1} a_{n1} - b_{n1} b_{n1}) ,$$

where

$$\rho_{ik} = \rho_i + \rho_k .$$

As a necessary condition for minimum, let

$$\frac{\partial \mathcal{L}}{\partial u_k} = 0, \qquad \frac{\partial \mathcal{L}}{\partial v_k} = 0, \qquad \frac{\partial \mathcal{L}}{\partial \psi_k^n[11']} = 0 ,$$

$$\frac{\partial \mathcal{L}}{\partial a_{nk}} = 0, \qquad \frac{\partial \mathcal{L}}{\partial b_{nk}} = 0, \qquad \frac{\partial \mathcal{L}}{\partial \phi_k^n[11']} = 0 .$$

The quantities ρ_i are not independent parameters. They are related with the boson amplitudes through the equation

$$\rho_1 = \frac{1}{2\Omega_1} \sum_{Ln2} (1-\rho_{12}) 2\Omega_L ((\psi_L^n[12])^2 + (\phi_L^n[12])^2 - 1) +$$

$$+ \frac{1}{2\Omega_1}(1 - 2\rho_1) \sum_n (a_{n1}^2 + b_{n1}^2 - 1) .$$

Neglecting the influence due to the pair vabrations (i.e., assuming that $a_{nk} = b_{nk} = 0$), we derive a system of equations which

1) connects the quasiparticle and phonon characteristics of the system;

2) turns into the known system of equations of the RPA at $\rho \to 0$. So, the secular equation has the form

$$1 = 2\chi_L \sum_{11'} \frac{(1-\rho_1-\rho_{1'})G_L^2[11'](\epsilon_1+\epsilon_{1'})}{(\epsilon_1+\epsilon_{1'})^2-(\omega_L^n)^2}, \tag{16}$$

where

$$G[11'] = \frac{1}{\sqrt{2}}(u_1 v_{1'} + v_1 u_{1'}) \frac{<1||Q_L||1'>}{\sqrt{2\Omega_L}},$$

$$\epsilon_s = [(E_s-\lambda)(u_s^2-v_s^2) + Gu_s v_s \sum_{s'} 2\Omega_{s'} u_{s'} v_{s'}(1-2\rho_{s'})], \tag{17}$$

The quantities u_k and v_k are determined by the formulae

$$u_k = \frac{1}{2}(1 + \frac{R_k}{\sqrt{r_k^2+R_k^2}}),$$

$$v_k = \frac{1}{2}(1 - \frac{R_k}{\sqrt{r_k^2+R_k^2}}), \tag{18}$$

where

$$R_k = (E_s-\lambda)(1-2\rho_k),$$

$$r_k = r_k^{(1)} + r_k^{(2)} = (1-2\rho_k)\frac{G}{2}\sum_1 2\Omega_1 u_1 v_1(1-2\rho_1) +$$

$$+\frac{1}{2}\sum_{Ln}\chi_L(\sum_{22'}<2||Q_L||2'>(u_2 v_{2'}+v_2 u_{2'})(1-\rho_{22'})(\psi_L^n[22']+\phi_L^n[22'])\times \tag{19}$$

$$\times\frac{1}{2\Omega_k}<k||Q||k>(1-\rho_{kk})(\psi_L^n[kk]+\phi_L^n[kk]).$$

It should be noted that the term $r_k^{(2)}$ including the influence of multipole forces to the value of r_k, appears when $\rho_i \neq 0$ only, i.e., when condition (13) is not fulfilled.

The boson amplitudes are equal to

$$\psi_L^n[11'] = \frac{N_L^n Q_L[11']}{(\epsilon_1+\epsilon_{1'}-\omega_L^n)},$$

$$\phi_L^n[11'] = \frac{N_L^n Q_L[11']}{(\epsilon_1+\epsilon_{1'}+\omega_L^n)}, \tag{20}$$

where

$$N_L^n = \chi_L \sum_{22'} (1-\rho_{22'}) Q_L[22'](\psi_L^n[22'] + \phi_L^n[22']).$$

Using the normalization condition

$$\sum_{11'} (1-\rho_{11'})((\psi_L^n[11'])^2 - (\phi_L^n[11'])^2] = 2$$

we get the following equation for calculating the quantity

$$(2N_L^n)^2 \omega_L^n \sum_{11'} \frac{(1-\rho_{11'})Q_L^2[11'](\epsilon_1+\epsilon_{1'})}{(\epsilon_1+\epsilon_{1'})^2-(\omega_L^n)^2} = 1 .$$

With this expression for N_L^n, one can write (18) as follows:

$$r_k = r_k^{(1)} + r_k^{(2)} = (1-2\rho_k)\frac{G}{2}\sum 2\Omega_1 u_1 v_1 (1-\rho_1) +$$

$$+ \frac{1}{\sqrt{2}} \sum_{Ln} \sqrt{2\Omega_L} N_L^n \frac{1}{2\Omega_k} \langle k\|Q_L\|k\rangle (1-\rho_{kk})(\psi_L^n[kk]+\phi_L^n[kk]). \tag{21}$$

Inclusion into the commutation relations (11) of the terms proportional to ρ_i causes appearance of ρ_i in the secular equation (16) and in the expression for quasiparticle energies ϵ_s (17).

The quasiparticle energies of the states with large ρ_s (for which the number of quasiparticles in the ground state is large) may differ essentially from the quantities calculated by the superfluid nuclear model [15].

This difference is the larger, the stronger are the collective properties of the lowest excitations. The contribution of these states into the secular equation (16) will be suppressed strongly. Following the results of ref. [10], which show that large ρ_i will have the states lying near the Fermi surface, one may conclude that the corrections can change essentially the structure of the lowest states. The quantity r_k (19),(21) in this consideration plays the role of the gap parameter. Two effects cause the change of r_k in comparison with the superfluid model[15]. First, due to the presence of ρ_k in the first term $r_k^{(1)}$ (19), r_k decreases in comparison with the superfluid model. Second, the second term $r_k^{(2)}$ will increase the value of r_k by its coherence. These effects are much stronger when the first excited state of an even-even nucleus is more collective. Taking into account the fact that the largest values of ρ_i have the states lying near the Fermi level, we can state that the strongest changes of the gap can be expected for the states in that region.

The change of the structure of the lowest excited states causes the change of the electromagnetic transition probability.

The reduced probability of electric transitions of multipolarity L is expressed by the formula

$$B(EL; L \to 0) =$$

$$= \{ \sum_{11'} (1-\rho_{11'}) Q_L [11'] (\psi_L^n [11'] + \phi_L^n [11'])\}^2 . \tag{22}$$

It is shown in refs. [12-14] that this value is less than that calculated in the RPA.

III. NUMERICAL RESULTS AND THEIR ANALYSIS

It is seen from the secular equation (16) that the problem stated can be solved in two stages. First, one should determine the values of ρ_i and then, having calculated all the rest unknown quantities (18),(19),(21), derive eq.(16). It is shown in ref. [12] that the values of ρ_i can be found by solving the linear system

$$(2\Omega_1 + f_1)\rho_1 + \sum_2 f_{12}\rho_2 = \rho_1 , \tag{23}$$

where

$$f_1 = \sum_2 f_{12} ; \quad f_{12} = \sum_{Ln} 2\Omega_L (\phi_L^n [12])^2 .$$

If the calculations use a large basis of single-particle states, the rank system (23) is large. One of its properties is that the diagonal terms exceed considerably the non-diagonal ones. This allows one to use the method from ref. [18], for the solution of system (23).

In formulae (19),(21) and (23) the summation over the phonon numbers is performed, and the phonons of different multipolarity are included in the sums.

Following the results of ref. [10] we have taken into account the first quadrupole and octupole phonons only. The influence of other high-lying collective states and of the spin-multipole phonons is insignificant for the quantity ρ_i [10].

The numerical calculations have been performed for $^{144-150}$Sm. The Sm isotopes change sharply collectivity of their 2_1^+ and 3_1^- states. The ^{144}Sm is a nucleus with one closed shell and its lowest excitations are not very collective. The ^{150}Sm is a good example of a transitional nucleus in which B(E2) amounts to dozens of single-particle units.

Table 1.

Energies and B(EL)-values for Sm isotopes. a) denotes the RPA results; b) the results taking into account the number of quasiparticles in the ground state; c) the results taking into account the correlations of quasiparticle and phonon excitations.

Nuclei	Index	$E(2_1^+)$ [MeV]		$B(E2)\uparrow$ e^2b^2		$E(3_1^-)$ [MeV]		$B(E3)\uparrow$ e^2b^3	
		theor.	exp.	theor.	exp.	theor.	exp.	theor.	exp.
^{144}Sm	a	2.15		0.254		2.30		0.147	
	b	2.17	1.66	0.234	0.25	2.34	1.810	0.136	-
	c	2.14		0.248		2.34		0.132	
^{146}Sm	a	1.174		0.482		2.0		0.193	
	b	1.478	0.747	0.293	-	2.23	1.381	0.144	-
	c	1.00		0.472		2.023		0.16	
^{148}Sm	a	0.87		0.679		1.65		0.247	
	b	1.45	0.550	0.309	0.70	2.09	1.162	0.156	0.25
	c	1.09		0.428		1.87		0.171	
^{150}Sm	a	0.49		1.33		1.50		0.262	
	b	1.43	0.334	0.361	1.37	2.05	1.071	0.151	0.31
	c	1.03		0.515		1.85		0.162	

2 Global analyses of nucleon EM structure data by the standard VMD model

There is a decomposition of nucleon electric $G_E(t)$ and magnetic $G_M(t)$ FF's into isoscalar and isovector parts of the Dirac and Pauli FF's as follows

$$\begin{aligned}
G_E^p(t) &= [F_{1N}^s(t) + F_{1N}^v(t)] + \frac{t}{4m_p^2}[F_{2N}^s(t) + F_{2N}^v(t)] \\
G_M^p(t) &= [F_{1N}^s(t) + F_{1N}^v(t)] + [F_{2N}^s(t) + F_{2N}^v(t)] \\
G_E^n(t) &= [F_{1N}^s(t) - F_{1N}^v(t)] + \frac{t}{4m_n^2}[F_{2N}^s(t) - F_{2N}^v(t)] \\
G_M^n(t) &= [F_{1N}^s(t) - F_{1N}^v(t)] + [F_{2N}^s(t) - F_{2N}^v(t)].
\end{aligned} \tag{1}$$

Cesselli, Nigro and Voci [1] have saturated the latter by means of the following five $\omega(782)$,$\phi(1019)$,$\phi'(1660)$,$\rho(773)$,$\rho'(1600)$ resonances considering the standard VMD model parametrization

$$\begin{aligned}
F_1^s(t) &= \sum_{s=\omega,\phi,\phi'} \frac{m_s^2}{m_s^2 - t}(f_{sNN}^{(1)}/f_s); \quad F_2^s(t) = \sum_{s=\omega,\phi,\phi'} \frac{m_s^2}{m_s^2 - t}(f_{sNN}^{(2)}/f_s); \\
F_1^v(t) &= \sum_{v=\rho,\rho'} \frac{m_v^2}{m_v^2 - t}(f_{vNN}^{(1)}/f_v); \quad F_2^v(t) = \sum_{v=\rho,\rho'} \frac{m_v^2}{m_v^2 - t}(f_{vNN}^{(2)}/f_v),
\end{aligned} \tag{2}$$

and carrying out a simultaneous analysis of 199 experimental points consisting of the proton and neutron space-like and the proton time-like data. In a fit of the proton FF data in the time-like region the $1/(m^2-t)$ factor has been modified to $1/(m^2-t+im\Gamma)$, where Γ is the width of unstable vector-mesons under consideration.

The main result of this analysis is an estimate of the $e^+e^- \to n\bar{n}$ cross-section to be about 100 times larger than the $e^+e^- \to p\bar{p}$ cross-section. At the same time, a strong constraint on the fit from the time-like data has been observed. Really, if in a fit only space-like data are used, the calculated value of $|G_E^p(4m_p^2)| = |G_M^p(4m_p^2)|$ at the proton-antiproton threshold is 2.27 against the measured value 0.51. This means that it is impossible to reproduce the correct time-like behaviour of the proton EM FF's in the framework of the standard VMD model on the basis of the space-like data only. The latter property is observed in all our further analyses as well.

We have used the same model analysing the nucleon FF data, however, extended for proton FF's up to $t = -33GeV^2$ and for electric and magnetic neutron FF's up to $t = -4GeV^2$ and $t = -10GeV^2$, respectively. We confirm roughly the Cesselli, Nigro, Voci [1] result for a rate of $\sigma_{tot}(e^+e^- \to n\bar{n})$ to $\sigma_{tot}(e^+e^- \to p\bar{p})$ but it is no more acceptable from the statistical point of view of an elaboration of experimental data as we have $\chi^2/NDF = 7025/382$ to be compared with $\chi^2/NDF = 359/189$ in ref.[1]. The obtained values of the corresponding coupling ratios are presented in Table 1, where they are compared also with the results of Cesselli, Nigro, Voci [1].

We have tried to analyse all nucleon FF data also by a more sophisticated standard VMD model of the nucleon EM structure taking into account the newest experimental

Table 2.

The quasiparticle energies ϵ_s and gap parameters $r_k^{(2)}$ and r_k for some states near Fermi surface in ^{150}Sm. (The gaps of superfluid model are $C_N \sim 1.29$ MeV, $C_Z \sim 1.22$ MeV).

nl_j neutrons	ϵ_s [MeV] superfl. model	ϵ_s [MeV] this paper	$r_k^{(2)}$ [MeV]	r_k [MeV]	nl_j protons	ϵ_s [MeV] superfl. model	ϵ_s [MeV] this paper	$r_k^{(2)}$ [MeV]	r_k [MeV]
$1h_{11/2}$	5.35	5.33	0.04	1.18	$2p_{1/2}$	6.86	6.85	0.0	1.12
$2d_{5/2}$	5.12	5.09	0.02	1.17	$1g_{9/2}$	6.28	6.26	0.02	1.13
$2f_{7/2}$	1.29	1.19	0.69	1.60	$1g_{7/2}$	1.44	1.36	0.33	1.36
$1h_{9/2}$	2.18	2.11	0.20	1.33	$2d_{5/2}$	1.26	1.18	0.42	1.36
$3p_{3/2}$	2.54	2.48	0.19	1.25	$1h_{11/2}$	1.75	1.68	0.31	1.35
$1i_{13/2}$	2.57	2.51	0.20	1.32	$3s_{1/2}$	2.54	2.50	0.	1.06
$3p_{1/2}$	3.61	3.58	0.0	1.15	$2d_{3/2}$	2.64	2.60	0.07	1.15

It is seen from table 2 that the parameter of the gap r_k changes differently for different states. As has been mentioned, r_k contains two terms $r_k^{(1)}$ and $r_k^{(2)}$ with opposite signs. It is seen from (19) and (21) that $r_k^{(2)}$ is directly connected with the constants of the multipole forces therefore it will be sensitive to the change of collectiveness of the 2^+ and 3^- states. The largest values of $r_k^{(2)}$ have the single-particle states near the Fermi surface. In spite of large ρ_i in these states (it is seen from (19) and (21) that ρ_i decrease $r_k^{(2)}$), large amplitudes ψ and ϕ (20) entering into (19) and (21), cause large values of $r_k^{(2)}$.

The term $r_k^{(1)}$ decreases the gap in contrast with the superfluid model. These differences are essential in ^{150}Sm , in which ρ_i are the largest. As it is seen from table 2, the values of r_k are lower than the gap parameters C_N and C_Z, calculated by the superfluid model. Only the single-particle states near the Fermi surface (neutron $2f_{7/2}$, $2p_{3/2}$, $1i_{13/2}$ and proton $1g_{7/2}$, $2d_{5/2}$, $1h_{11/2}$) have a larger gap than in the superfluid model.

New gaps r_k cause new u, v -coefficients and change the one-quasiparticle energies ϵ_s . It follows from (17) that ϵ_s should be less than the quasiparticle energies of the superfluid model. However, the increasing gap compensates this decrease and ϵ_s is only somewhat lower than those values calculated by the superfluid model. A slight but essential shift of the first poles of eq.(16) occurs. At that the multipole force constants κ_L do not change, and the change of ϵ_s and u, v -coefficients results in increasing collectiveness of the lowest 2^+ and 3^- states. Therefore, in table 1c) the energies of the 2^+_1 and 3^-_1 states are lower than in b). The electric transition probabilities are also larger in c) than in b).

By comparing the rows a) and c) of table 1, one can see that the number of quasiparticles in the ground state and the correlations of phonon and quasiparticle excitations lead to two opposite effects giving almost the same corrections. In the semimagic ^{144}Sm and neighbouring ^{146}Sm, in which collectiveness of the first 2^+ and 3^- states is weaker, these corrections are small and compensate each other completely. With increasing collectiveness of the 2^+_1 and 3^-_1 states, the number of quasiparticles in the ground state of even nuclei increases, the values of the vector ρ become larger and the first effect predominates. This changes essentially the energies of the 2^+_1 and 3^-_1 states and B(E2)- and B(E3)-values in ^{148}Sm and ^{150}Sm.

CONCLUSION

The above investigation shows that there exist correlations connecting the amplitudes of phonon and quasiparticle excitations. They cause the change of the gap of quasiparticle states. Its value depends to a great extent on the multipole force constants. On the other hand, the near gap changes the quasiparticle energies, Bogolubov's coefficients and consequently, the collectiveness of the lowest excitations.

The numerical calculations of the Sm isotopes have shown that influence of the multipole forces on the gap is notable for the states near the fermi level. This effect, however, is compensated to a great extent by the influence of quasiparticles in the ground state. As a result the gap and the quasiparticle energies of states only slightly differ from the values calculated within the superfluid nuclear model.

The influence of the number of quasiparticles in the ground state on the 2^+_1 and 3^-_1 states diminishes their collectiveness. However, the collectiveness of the 2^+_1 and 3^-_1 states increases, if one takes into account the correlations between the number of quasiparticles in the ground state and the phonon amplitudes ψ and ϕ.

The calculations have shown that the effects under consideration are small and compensate each other completely in the semimagic ^{144}Sm and neighbouring ^{146}Sm.

In the transitional nuclei the number of quasiparticles in the ground state is so large, that its inclusion causes a complete loss of collective properties of the first 2^+ and 3^- states. In these nuclei to abequatly describe the properties of collective states, one should evidently include into the Hamiltonian the terms taking into account more complex relations of quasiparticle degrees of freedom as has been mentioned in ref. [14].

The numerical calculations allow one to conclude that it is reasonable to use the RPA as an initial approximation, since the effects under consideration are small in semimagic and adjacent nuclei. Therefore, the RPA-phonons used as a basis in the quasiparticle-phonon nuclear model[1] are a good basis. However, it is impossible to use the RPA-phonons as a basis for the description of the fragmentation of simple excitations in the nuclei, where the first excited states are strongly collective (as in ^{150}Sm). Those nuclei require additional investigations of the role of other terms in the Hamiltonian[4], which form the first 2^+ and 3^- excited states.

The authors are grateful to prof. V.G.Soloviev for interest in this work and valuable remarks. They thank also A.I.Vdovin and R.V.Jolos for useful discussions.

REFERENCES.

1. Soloviev V.G. Sov.Particles and Nuclei, 1978, 9, p.810; Nucleonica, 1979, 23, p.1149.
2. Lin K.F., Brown G.R. Nucl.Phys., 1976, A265, p.385.
3. Pyatov N.I., Baznat M.I. Yad.Fiz., 1979, 30, p.1219. Kamerdzhiev S.P. Electromagnetic Interactions of Nuclei at Low and Medium Energies, Nauka, 1979, p.93.
4. Li C.T. et al. Phys.Rev.C, 1979, 19, p.1002. Mitroshin V.E. Izv.Akad. Nauk SSSR (ser.fiz.), 1976, 40, p.126.
5. Soloviev V.G., Stoyanov Ch., Vdovin A.I. Nucl.Phys., 1977, A288, p.376.
6. Soloviev V.G., Stoyanov Ch., Voronov V.V. Nucl.Phys., 1978, A304, p.503.

7. Soloviev V.G., Stoyanov Ch., Voronov V.V. Phys.Lett., 1978, 79B, p.187.
Vdovin A.I. et al. Yad.Fiz., 1979, 30, p.923.
8. Ponomarev V.Yu. et al. Nucl.Phys., 1979, A323, p.446.
9. Malov L.A., Soloviev V.G. Sov.Particles and Nuclei, 1980, 11, p.301.
Soloviev V.G., Vdovin A.I. Proc. of IV EPS Balaton Conf. on Nucl.Phys. ed by A.Kiss, J.Nemeth, J.Zimamyi, 1979, v.1, p.131.
10. Soloviev V.G., Stoyanova O., Stoyanov Ch. JINR, E4-80-75, Dubna, 1980.
11. Nesterenko V.O., Soloviev V.G., Halkin V. JINR, E4-80-143, Dubna, 1980.
12. Hara K. Prog.Theor.Phys., 1964, 32, p.88.
13. Ikeda K., Udagawa T., Yamamura H. Prog.Theor.Phys., 1965, 33, p.22.
14. Jolos R.V., Rybarska V. JINR, E4-5578, Dubna, 1971.
15. Soloviev V.G. Theory of Complex Nuclei, Oxford, Pergamon Press., 1976.
16. Stoyanov Ch. Theor.Mat.Fiz., 1979, 40, p.422.