сообщения
объединенного
института
ядерных
исследований
дубна

E4-80-132

V.G.Soloviev, Ch.Stoyanov, A.I.Vdovin

THE EFFECT OF DIFFERENT TYPES OF PHONONS AND COMPLEX CONFIGURATIONS ON THE DISTRIBUTION OF ONE-QUASIPARTICLE STRENGTH

1980

In recent years the investigation of highly excited nuclear states have raised much interest of experimenters and theoreticians. The reason is the discovery of new giant resonances [1] and resonance-like structures in the cross sections of the one-nucleon transfer reactions [2]. The main characteristics of the giant resonances and the resonance-like structures are explained by the fragmentation of the simplest nuclear excitations (one-phonon or one-quasiparticle) over more complex nuclear states. The fragmentation is extensively studied in the framework of the semimicroscopic quasiparticle-phonon nuclear model [3]. This model gives the consistent microscopic description of the spreading width of giant resonances, the E1- and M1-radiative strength functions and neutron strength functions in even-even spherical nuclei [4]. In the odd-mass spherical nuclei the s- and p-wave neutron strength functions and deeply-bound hole states have been investigated [5,6]. In refs. [5,6] the model wave function of an odd-mass nucleus has the form:

$$\Psi_{\nu}(JM) = C_{J\nu}\{\alpha^{+}_{JM} + \sum D^{\lambda i}_{j}(J\nu)[\alpha^{+}_{jm}Q^{+}_{\lambda\mu i}]_{JM}\}\Psi_{0}, \tag{1}$$

i.e., the interaction of one-quasiparticle states with "quasiparticle + phonon" states only has been taken into account. In formula (1) α^{+}_{jm} is the quasiparticle creation operator with shell quantum numbers $(n, \ell, j) \equiv j$ and momentum projection m; $Q^{+}_{\lambda\mu i}$ is the phonon creation operator with momentum λ, its projection μ and root number i. There exists a set of phonons with given λ having different excitation energies. The index i distinguishes these different phonons. Ψ_0 is the ground state wave function of the neighbouring even-even nucleus. The calculations with a more complex wave function than (1) are very difficult. A numerical method for solving the equations of the quasiparticle-phonon model for the wave function with "quasiparticle + two phonons" components has recently been suggested [7].

In the present paper we investigate the effect of "quasiparticle + two phonons" components on the fragmentation of the one-quasiparticle strength at intermediate and high excitation energies. We investigate also the role of the phonons of different multipolarities for the description of the fragmentation. The model wave function has the following form:

$$\Psi_\nu(JM) = C_{J\nu}\{\alpha^+_{JM} + \sum_{\lambda i j} D^{\lambda i}_{j}(J\nu)[\alpha^+_{jm} Q^+_{\lambda\mu i}]_{JM} +$$

$$+ \sum_{\substack{\lambda_1 i_1 \lambda_2 i_2 \\ jI}} F^{\lambda_1 i_1 \lambda_2 i_2}_{jI}(J\nu)[\alpha^+_{jm}[Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2}]_{IM'}]_{JM}\}\Psi_0 . \qquad (2)$$

We take into account a large number of phonons with $\lambda^\pi = 1^\pm, 2^\pm, 3^\pm ... 7^\pm$. The derivation of the model equations for the coefficients $C_{J\nu}$, $D^{\lambda i}_{j}$ and $F^{\lambda_1 i_1 \lambda_2 i_2}_{jI}$ and energy $\eta_{J\nu}$ of the state (2) is given in refs. /8,9/. In refs. /7,9/ the numerical iteration method for solving these equations has been described in detail. In the present calculation the method of strength function /10,11/ is used. The strength function $C^2(\eta)$ has been defined /10/ as follows:

$$C^2(\eta) = \frac{\Delta}{2\pi} \sum_\nu \frac{1}{(\eta-\eta_{J\nu})^2 + \Delta^2/4} C^2_{J\nu} . \qquad (3)$$

The function $C^2(\eta)$ describes the change of the average value of the square of one-quasiparticle components $C^2_{J\nu}$ with excitation energy η. The averaging of $C^2_{J\nu}$ is performed over the energy interval Δ which should be much larger than the mean distance between neighbouring roots $\eta_{J\nu}$. A particular form of $C^2(\eta)$ corresponding to the wave function (2) has been given in ref. /9/ and to the wave function (1) in refs. /5,6,10/.

At first we discuss the model parameters. The parameters of the quasiparticle-phonon model are the single-particle potential parameters and the constants of the residual effective interaction. The justification and choice of the parameters and the analysis of their influence on the phonon properties can be found in refs. /4,5/. The parameters of the Saxon-Woods single-particle potential are chosen as in ref. /12/. For the constants of separable multipole forces with $\lambda > 3$ and the spin-multipole forces the estimates of ref. /13/ are used. The constants of the pairing interaction, separable dipole, quadrupole and octupole forces are determined by the experimental data. The constants of quadrupole and octupole forces are determined so as to describe correctly the 2^+_1 and 3^-_1 energies in the RPA, the $B(E2, 0^+ \to 2^+_1)$ and $B(E3, 0^+ \to 3^-_1)$ values being somewhat overestimated in comparison with experiment. The interaction between quasiparticles and phonons has no any additional free parameters. It depends on the phonon structure and single-particle matrix elements of the residual forces /8,9/.

The parameter Δ, which has appeared in the definition of the strength function (3) is chosen to be equal to 0.2 MeV in ^{57}Ni and 0.5 MeV in ^{119}Sn and ^{123}Te. This value is the standard resolution in the one-nucleon transfer experiments at the excitation energy $E_x > 5 \div 10$ MeV and the accuracy of the quasiparticle-phonon model in describing the experimental data. The Δ-value should be of the same order of magnitude at the energies $E_x \gtrsim 5 \div 10$ MeV if the "smearing" of $C_{J\nu}$ is considered as a way of taking into account the influence of the single-particle continuum /14,15/ and of complex configurations /6,15,16/.

It is considered that the fragmentation of simplest states even at high excitation energy is determined mainly by the interaction with the low-lying vibrations. The dimension of the basis formed by the quadrupole and octupole phonons is not large. Let us consider the density of states of the type of "quasiparticle plus phonon" and "quasiparticle plus two phonons" in a spherical nucleus. Figure 1 shows the histogram of a number of states $N(E_x)$ with $J^\pi = 9/2^+$ in the energy interval $\Delta E_x = 1$ MeV as a function of the excitation energy in ^{123}Te. The calculations have been performed with two phonon bases. The comparison of different histograms in figs. 1a and 1b shows that the limitation of the phonon basis by the quadrupole and octupole phonons decreases the state density by

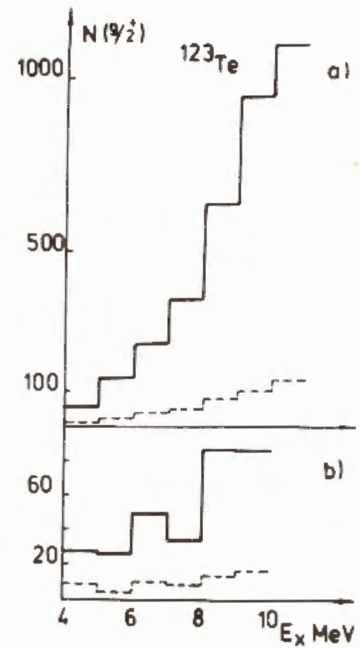

Fig. 1. Histogram of a number of states with $J^\pi = 9/2^+$ in the interval $\Delta E_x = 1$ MeV as a function of the excitation energy E_x in ^{123}Te. a) Histogram of a number of states of the type $[\alpha^+ Q^+ Q^+]_{9/2^+}$, the solid line is the calculation with a full phonon basis, the dashed line is the calculation with the quadrupole and octupole phonons only; b) Histogram of a number of states of the type $[\alpha^+ Q^+]_{9/2^+}$, the solid line is the calculation with a full phonon basis, the dashed line is the calculation with the quadrupole and octupole phonons only.

five times. This difference is increasing with E_x. Such a decrease in the state density simplifies greatly the calculations. Therefore Doll[15], Koeling and Iachello[16] have described the fragmentation of deeply-bound hole states in terms of a doorway-state-like picture, in which the hole strength is first fragmented due to the strong coupling to the lowest quadrupole and octupole phonons, the fragments being subsequently spread over many underlying non-collective states in a statistical way. The Lorentz or Gauss form of the normalized weight function has been used for folding the stectroscopic factor for the collective doorway states. The width of the weight function has been calculated by the formula

$$\Delta(E_x) = 2\pi N(E_x) \langle V^2 \rangle ,$$

where $\langle V^2 \rangle$ is the average square matrix element of the interaction with the noncollective states of a more complex structure. Such a semi-statistical approach is partly justified by the recent results of refs.[6,11]. It has been shown that the interaction with the phonons of high multipolarity and spin-multipole phonons results in the broadening of peaks and smoothing of fluctuations of the strength function. But the gross-structure of the function $C^2(\eta)$ is caused by the interaction with the low-lying quadrupole and octupole vibrations.

However, this picture sometimes is invalid. Among the phonons at intermediate and high energies there are the collective ones whose interaction with the one-quasiparticle state will change the strength function essentially like the interaction with the low-lying quadrupole and octupole vibrations. Most of these high-lying collective phonons are at the energies $E_x > 10$ MeV but some of them are at lower energy, for instance the low-lying octupole resonance (LEOR)[4,17]. We show the influence of the LEOR on the fragmentation of the $1g_{9/2}$ state in the tin isotopes. In our calculations the LEOR in the Sn isotopes is exhibited as a collective 3^--state with the energy $\omega \simeq 5$ MeV and $B(E3) = 0.25 B(E3, 0^+ \to 3^-_1)$[4]. Figure 2 shows the strength function of the neutron hole state $1g_{9/2}$ in ^{119}Sn calculated with the wave function (1) for different phonon bases. The wave function of the simplest structure (it includes six components) corresponds to fig.2c. The interaction of quasiparticles with the first quadrupole and octupole phonons only is taken into account. A bulk of the $1g_{9/2}$-strength in this case is distributed between three states. The main peak of the function $C^2(\eta)$ splits when the interaction with the quadrupole and octupole high energy phonons is included (fig.2b). This is caused by the interaction with the LEOR. The excitation energy ϵ_J of the $1g_{9/2}$ hole state in lighter tin isotopes is less than in ^{119}Sn and the interaction with the LEOR in these nuclei changes the function $C^2(\eta)$ at the energies above the main peak. Figure 2a shows that the inclusion of phonons of other multipolarities enhances the splitting of the main peak of the function $C^2(\eta)$ and causes its further broadening. Thus, the high-lying collective states strongly influence the fragmentation of hole states and the interaction with them cannot be taken into account by introducing the doorway-state effective width only.

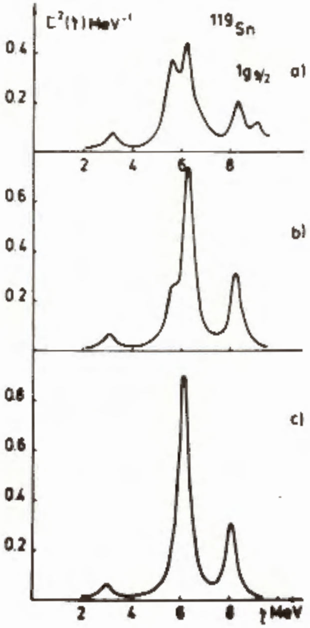

Fig. 2. Strength function of the $1g_{9/2}$ neutron hole state in ^{119}Sn. The calculation with the wave function (1). a) the calculation with a full phonon basis; b) the calculation with the quadrupole and octupole phonons only, having the energy $\omega_{\lambda i} \leq 11$ MeV; c) the calculation with the first quadrupole and octupole phonons.

Let us consider the influence if the components of the $\alpha^+ Q^+ Q^+$-type on the fragmentation of hole states. Their importance is determined by the strength of the quasiparticle-phonon interaction and the density of $\alpha^+ Q^+ Q^+$-states. The comparison of figs. 1a and 1b shows that the density of $\alpha^+ Q^+ Q^+$-states at the energies $E_x \simeq 6 \div 8$ MeV is 3-5 times as large as the density of states of the $\alpha^+ Q^+$-type, and with increasing energy E_x the difference between the densities increases. Figures 3c and 4c show the strength function $C^2(\eta)$ of the $1g_{9/2}$ neutron hole state in ^{119}Sn and ^{123}Te. The calculation has been performed with the wave function (2) and the full phonon basis. For comparison figs. 3b and 4b show the function $C^2(\eta)$ calculated taking into account the quadrupole and octupole phonons only, and figs. 3a and 4a show the strength function calculated with the wave function (1). As is seen from the figures, the interaction with the states $\alpha^+ Q^+ Q^+$ causes a stronger fragmentation of the $1g_{9/2}$ hole state and changes essentially the strength function, especially at the energies $E_x > 7$ MeV. The deep minimum of the function $C^2(\eta)$ at $\eta = 7.5$ MeV disappears

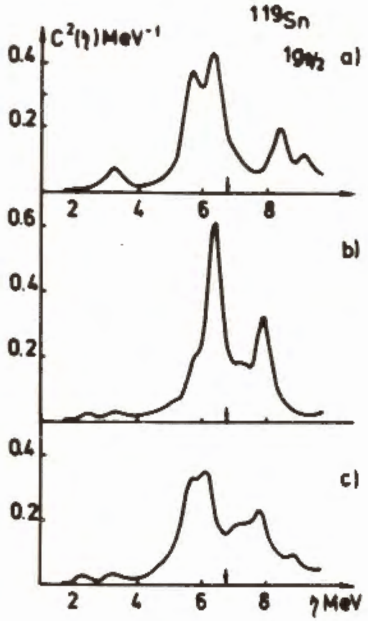

Fig.3. Strength function of the $1g_{9/2}$ neutron hole state in ^{119}Sn: a) the calculation with the wave function (1) and a full phonon basis; b) the calculation with the wave function (2). The quadrupole and octupole phonons only with the energy $\omega_{\lambda i} \leq 11$ MeV have been taken into account; c) the calculation with the wave function (2) and a full phonon basis. The arrow shows the position of the $1g_{9/2}$ one-quasiparticle level.

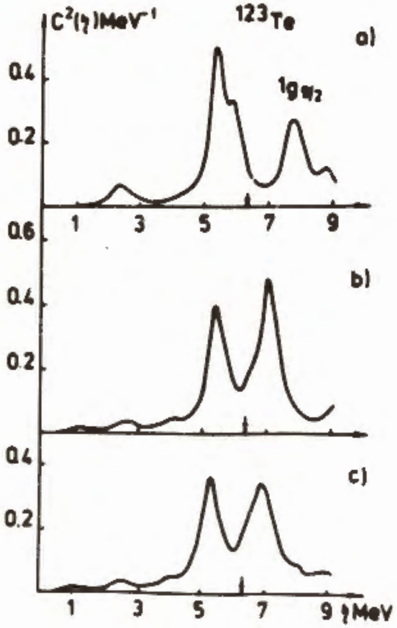

Fig.4. Strength function of the $1g_{9/2}$ neutron hole state in ^{123}Te: a) the calculation with the wave function (1) and a full phonon basis; b) the calculation with the wave function (2). The quadrupole and octupole phonons only with the energy $\omega_{\lambda i} \leq 11$ MeV have been taken into account; c) the calculation with the wave function (2) and a full phonon basis. The arrow shows the position of the $1g_{9/2}$ one-quasiparticle level.

in ^{119}Sn. In ^{123}Te the height of the main peak sharply diminishes, and the second peak, at an energy of $\eta = 7.5$ MeV, increases and shifts towards the lower energies. As a result both the peaks become of the same height and width. Most of these changes occur due to the interaction with the quadrupole and octupole phonons (see figs. 3b and 4b).

The influence of the $\alpha^+Q^+Q^\pm$-states on the fragmentation of the $1g_{9/2}$ hole state in ^{123}Te is stronger than in ^{119}Sn. The reason is that the first quadrupole and octupole phonons in the Te isotopes are more collective than in the tin isotopes. The higher the collectiveness of the phonon, the stronger it interacts with the quasiparticles. The decrease in the constants of the quadrupole and octupole forces increases the energies of the 2^+_1 and 3^-_1-levels and diminishes their collectiveness. The quasiparticle-phonon interaction weakens and the role of the $\alpha^+Q^+Q^+$-configurations decreases. In ref. [6] the values of quadrupole and octupole constants have been less than in the present paper and consequently the energies of the 2^+_1 and 3^-_1-states in even-even tin isotopes have been somewhat higher and their B(E2)-and B(E3)-factors lower than the experimental ones. At those values of constants of the effective forces the influence of the $\alpha^+Q^+Q^+$-configurations on the strength function is negligible.

Due to a rapid increase in "quasiparticle plus two phonons" state density with E_x (see fig.1a), their influence may turn to be strong at the energies $E_x \geq 15 \div 20$ MeV. As an example of a strong influence of the $\alpha^+Q^+Q^+$-states fig. 5 shows the strength distribution of the $1p_{3/2}$ neutron hole state with $\epsilon_J = 18.2$ MeV in ^{57}Ni. The strength function in fig. 5a is calculated with the wave function (2) and in fig. 5b with the wave function (1). The comparison of figs. 5a and 5b shows that the interaction with the $\alpha^+Q^+Q^+$-states distributes the $1p_{3/2}$ -strength (concentrated in the peak at $\eta = 22$ MeV) over the energy interval $\Delta\eta \simeq 4$ MeV at lower excitation energy ($\bar{\eta} \simeq 18$ MeV).

The maximum of the function $C^2(\eta)$ is as a rule shifted with respect to the unperturbed position of the hole level ϵ_J. In the Sn and Te isotopes this shift is of several hundred keV. But the centroid $\bar{E}_x$ of the function $C^2(\eta)$ calculated in the interval exhausting the most part of the hole state strength is close to ϵ_J. In the Sn and Te isotopes the value of $|\epsilon_J - \bar{E}_x|$ is 20-30 keV. This means that the parameters of the single-particle potential can be fitted using the centroids of the experimental distributions of single-particle strength as in ref. [18].

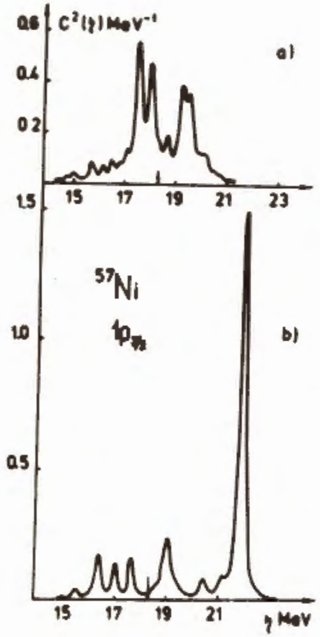

Fig.5. Strength function of the $1p_{3/2}$ neutron hole state in ^{57}Ni a) the calculation with the wave function (2) and a full phonon basis; b) the calculation with the wave function (1). The arrow shows the position of the $1p_{3/2}$ one-quasiparticle level.

The results of the present calculations have shown a large influence of the "quasiparticle + two phonons" components on the fragmentation of one-quasiparticle strength at intermediate and high excitation energies. The use of a large phonon basis is also necessary. Both factors should be taken into account for a correct description of the fragmentation of deeply bound hole states. The interaction of the high-lying one-quasiparticle state with the high-lying collective phonons is as important as the interaction with the low-lying quadrupole and octupole vibrations. Its influence cannot be taken into account by introducing the doorway-state effective width. The further investigation of the fragmentation of deeply-bound hole states and the detailed comparison with the available experimental data will be given in the next publication.

REFERENCES

1. Ritthan R., Walcher Th. Phys.Lett., 1971, 36B, p.563; Z.Naturforsch., 1972, 27A, p.1683; Satchler G.R. Phys. Rep., 1974, 14C, p.97; Pitthan R. Invited Seminar given at the 11th Masurian Summer School in Nuclear Physics, Mikolajki, Poland, 1978; Richter A. Contributed paper given at the IV Seminar on "Electromagnetic Interactions of Nuclei at Low and Medium Energies, IKDA, 78/3.
2. Sakai M., Kubo K.I. Nucl.Phys., 1972, A185, p.217. Siemssen R.H. Proc. of Int. Symp. on Selected Topics in Nucl. Struct., 1976, v.2, p.106. JINR, D-9920, Dubna, 1976; Tanaka M. et al. Phys.Lett., 1978, 78B, p.221;

Gerlic E. et al. Preprint Institut de Physique Nucleaire, Orsay, IPNO-PhN-79-20, 1979.
3. Soloviev V.G. Theory of Complex Nuclei. Pergamon Press, Oxford, 1976; Soloviev V.G. Particles and Nuclei, 1978, 9, p.810; Soloviev V.G. Nucleonika, 1979, 23, p.1149.
4. Soloviev V.G., Stoyanov Ch., Vdovin A.I. Nucl.Phys., 1977, A288, p.376; Soloviev V.G., Stoyanov Ch., Voronov V.V. Nucl.Phys., 1978, A304, p.503; Phys.Lett., 1978, 79B, p.187; Ponomarev V.Ju. et al. Yad.Fiz., 1979, 30, p.923; Nucl.Phys., 1979, A323, p.446; Soloviev V.G., Stoyanova O., Voronov V.V. JINR, E4-12500, Dubna, 1979.
5. Dambasuren D. et al. J.Phys.G: Nucl.Phys., 1976, 2, p.25.
6. Vdovin A.I., Stoyanov Ch., Chan Zuy Khuong. JINR, E4-12012, Dubna, 1978; Izv.Akad.Nauk SSSR (ser.fiz.), 1979, 43, p.999.
7. Stoyanov Ch. Teor.Mat.Fyz., 1979, 40, p.422.
8. Soloviev V.G., Vdovin A.I. Teor.Mat.Fyz., 1974, 19, p.275.
9. Soloviev V.G., Stoyanov Ch., Vdovin A.I. JINR, P4-12992, Dubna, 1979.
10. Bohr A., Mottelson B. Nuclear Structure, v.1, W.A.Benjamin Inc., New York, Amsterdam, 1969; Ochirbat G. JINR, P4-8774, Dubna, 1975.
11. Malov L.A., Soloviev V.G. Nucl.Phys., 1976, A270, p.87.
12. Chepurnov V.A. Yad.Fiz., 1967, 6, p.955; Takeuchi K., Moldauer P.A. Phys.Lett., 1969, 28B, p.384.
13. Bohr A., Mottelson B. Nuclear Structure, v.2. W.A.Benjamin Inc., New York, Amsterdam, 1974; Castel B., Hamamoto I. Phys.Lett., 1976, 65B, p.27.
14. Shlomo S., Bertsch G. Nucl.Phys., 1975, A243, p.507.
15. Doll P. Nucl.Phys., 1977, A292, p.165.
16. Koeling T. Iachello F. Nucl.Phys., 1978, A295, p.45.
17. Moss T.M. et al. Phys.Rev., 1978, C18, p.741.
18. Millener D.J., Hodgson P.E. Nucl.Phys., 1973, A209, p.59; Bear K., Hodgson P.E. J.Phys.G: Nucl.Phys., 1978, 4, p.L287.

Received by Publishing Department
on February 18 1980.