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E4-82-251

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FRAGMENTATION OF NEUTRON HOLE STATES IN $^{111,115}Sn$

Submitted to "Ядерная физика"

1982

INTRODUCTION

The first experimental data on the excitation of deep hole nuclear shells in the one-nucleon transfer reactions have emerged in the early seventies. At present the investigations of these resonance-like structures are in progress and an extensive experimental material is available (see, for instance, the review of S.Gales [1]). The study of deep hole states has been stimulated by the use of polarized particle beams, that allowed one to reliably determine the spin and parity of excited levels [2-5]. In the Sn isotope, for instance, alongside with a more accurate information on a rather well studied subshell $1g_{9/2}$ there have been obtained quantitative data on the strength distribution of deeper-lying hole $2p$ -subshells [4,5].

The fragmentation of deep hole states in medium and heavy nuclei has been investigated theoretically in papers [6-8]. In the last two papers [7,8] the quasiparticle-phonon nuclear model [9] has been used. According to the calculations the interaction of hole excitations with the quadrupole and octupole vibrations of the core plays the dominating role in the formation of the gross-structure of the hole state strength function. It has been shown in paper [8] that for a quantitative description of the available experimental data, one should take into account the interaction with a large number of states of the type "quasiparticle plus phonon" and "quasiparticle plus two phonons"; of particular importance is the coupling with the collective phonon excitations of intermediate energy (e.g., with the low-lying octupole resonance). A consistent consideration of these effects allowed one to get a satisfactory agreement with the experimental data.

The increasing availability of a new, more detailed and reliable experimental information prompts us to continue the theoretical study of the fragmentation of deep hole states, which has been undertaken in papers [7,8]. In the present paper we shall study the fragmentation of hole states $2p_{1/2}$, $2p_{3/2}$, $1f_{5/2}$ and $1f_{7/2}$ in the $^{111,115}Sn$ isotopes and of the state $1g_{9/2}$ in ^{111}Sn, which has not been studied in paper [8].

1. BASIC FORMULAE

The basic assumptions of the quasiparticle-phonon nuclear model are presented in paper [9]. A considerable part of the results obtained within this model is reviewed in [10]. In the present paper we shall follow a formalism of the model developed for the description of odd-A spherical nuclei. A more detailed presentation of this formalism is given in refs. [8,11]. The fragmentation of hole states is a result of their interaction with the phonon nuclear excitations. The contribution of one- and two-phonon excitations to this interaction has been taken into account in the present calculations, in other words the excited state wave function of an odd-A nucleus, written through the quasiparticle (α^+_{jm}) and phonon ($Q^+_{\lambda\mu i}$) creation operators, has the form

$$\Psi_\nu(JM) = C_{J\nu}\{\alpha^+_{JM} + \sum_{\lambda i j} D_j^{\lambda i}(J\nu)[\alpha^+_{jm}Q^+_{\lambda\mu i}]_{JM} + \sum_{\substack{\lambda_1\lambda_2 I\\ i_1 i_2 j}} F_{Ij}^{\lambda_1 i_1\lambda_2 i_2}(J\nu)[\alpha^+_{jm}[Q^+_{\lambda_1\mu_1 i_1}Q^+_{\lambda_2\mu_2 i_2}]_{I\mathfrak{M}}]_{JM}\}\Psi_0 \quad (1)$$

Ψ_0 is the ground state wave function of a neighbouring doubly even nucleus. The energies and structure of the one-phonon states, which are a superposition of the two-quasiparticle excitations, are calculated in the RPA. The phonons of natural parity are generated by separable multipole forces; and the phonons of unnatural parity, by separable spin-multipole forces. The one-phonon excitations of a doubly even nucleus unify both the lowest collective states (e.g., the 2^+_1, 3^-_1) and the collective states at intermediate and high excitation energy (various giant resonances: E1, E2, E3, M2, etc.). Most of the one-phonon states are noncollective, i.e., they are contributed by one or two two-quasiparticle components. The operator describing the quasiparticle-phonon interaction is

$$H_{qph} = -\frac{1}{2\sqrt{2}}\sum_{\lambda\mu i}(Q^+_{\lambda\mu i} + (-)^{\lambda-\mu}Q_{\lambda-\mu i})\sum_{j_1 j_2 \tau}\frac{f^{(\lambda)}_{j_1j_2}v^{(-)}_{j_1j_2}}{\sqrt{\mathcal{Y}^{\lambda i}_\tau}}B(j_1 j_2\lambda-\mu) + \text{h.c.} \quad (2)$$

$$B(j_1 j_2\lambda-\mu) = \sum_{m_1 m_2}(-)^{j_2+m_2}(j_1m_1j_2m_2|\lambda-\mu)\alpha^+_{j_1m_1}\alpha_{j_2-m_2};$$

$$v^{(-)}_{j_1j_2} = u_{j_1}u_{j_2} - v_{j_1}v_{j_2}.$$

For simplicity in (2) we have written just the part of the interaction which is related with the multipole phonons. We have used the following notation: $f^{(\lambda)}_{j_1j_2}$ is the reduced single-particle matrix element of the multipole operator; u_j and v_j are the Bogolubov transformation coefficients; $\mathcal{Y}^{\lambda i}_\tau$ is defined by the phonon structure (the expression for it is given in ref. [8]). The isotopic index τ acquires two values (n,p) and indicates which of the single-particle spectrum, neutron or proton, the single-particle states with quantum numbers j_1, j_2* belong to. Using the variational principle and taking into account the normalization of the wave function (1), in ref. [11] the system of equations has been obtained for the coefficients $D_j^{\lambda i}(J\nu)$ and energy of the state (1) $\eta_{J\nu}$:

$$\epsilon_J - \eta_{J\nu} - \frac{1}{\sqrt{2}}\sum_{\lambda i j}\Gamma(Jj\lambda i)D_j^{\lambda i}(J\nu) = 0,$$

$$\sum_{\lambda i j}D_j^{\lambda i}(J\nu)\Big[\Big(\epsilon_{j_1}+\omega_{\lambda_1 i_1} - \frac{1}{2}\sum_{\lambda_2 i_2 j_2}\frac{\Gamma^2(jj_2\lambda_2 i_2)}{\epsilon_{j_2}+\omega_{\lambda i}+\omega_{\lambda_2 i_2}-\eta_{J\nu}}\Big)\delta_{\lambda\lambda_1}\delta_{ii_1}\delta_{jj_1} +$$

$$+\frac{1}{2}\sum_{j_3}\frac{\Gamma(j_3j_1\lambda i)\Gamma(j_3j\lambda_1 i_1)}{\epsilon_{j_3}+\omega_{\lambda_1 i_1}+\omega_{\lambda i}-\eta_{J\nu}}[(2j_1+1)(2j_2+1)]^{1/2}\begin{Bmatrix}\lambda & j_3 & j_1\\ \lambda_1 & J & j\end{Bmatrix}\Big] = \frac{1}{\sqrt{2}}\Gamma(Jj_1\lambda_1 i_1), \quad (3)$$

where ϵ_j and $\omega_{\lambda i}$ are the energies of one-quasiparticle and one-phonon states.

$$\Gamma(Jj\lambda i) = \left(\frac{2\lambda+1}{2J+1}\right)^{1/2}\frac{f^{(\lambda)}_{Jj}v^{(-)}_{Jj}}{\sqrt{\mathcal{Y}^{\lambda i}_\tau}} \sim \langle\Psi_0\alpha_{JM}|H_{qph}|[\alpha^+_{jm}Q^+_{\lambda\mu i}]_{JM}\Psi_0\rangle$$

The system of equations (3) is obtained as a result of some approximations. In particular, the fermion structure of phonons has been neglected, i.e., we assumed $[\alpha^+_{jm}, Q^+_{\lambda\mu i}]=0$. We did not take into account the matrix elements of type

$$\langle\Psi_0\alpha_{JM}|H_{qph}|[\alpha^+_{jm}[Q^+_{\lambda_1\mu_1 i_1}Q^+_{\lambda_2\mu_2 i_2}]_{I\mathfrak{M}}]_{JM}\Psi_0\rangle$$

that means the neglect of the effects of anharmonicity of the core excitations. Besides, a specific structure of the wave function (1) does not allow one to take into account fully the correlations in the ground state of an odd-A nucleus caused by the quasiparticle-phonon interaction. Some of these effects have been investigated earlier [12], and it may be expected that they do not influence the fragmentation of the one-quasiparticle component considerably.

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*Here j implies three numbers $j \equiv n\ell j$.

The numerical solution of the system (1) is a very complicated problem from the computational viewpoint, even with the admissible limitations to the dimension of the single-particle and phonon bases. Therefore, as in papers [7,8], we shall not calculate the values $C_{J\nu}$ for each state $\Psi_\nu(JM)$, but the strength function $C^2(\eta)$ describing the dependence of the averaged over the interval Δ coefficient $C^2_{J\nu}$ on energy η

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

The calculation of the strength function turns out to be a more simple problem due to some specific properties of the system of equations (3). The method of strength functions has been proposed in paper [13]; here, one can find a thorough presentation and treatment of this method. The choice of the parameter Δ depends on the formulation of the problem. In paper [14] it has been analysed how to choose the value of Δ to avoid additional errors in the results of calculation, caused by an artificial broadening of states. The dependence of the results of averaging on the form of the weight function was also investigated in this paper.

In this present paper the parameters of the model Hamiltonian are taken the same as in paper [8]. The value of $\Delta =$ $= 0.2$ MeV. This value is considerably smaller than that in refs. [7,8] $\Delta = 0.5$ MeV. The reason for the decrease of Δ in the present calculations is that the experimental information obtained recently [5] is very detailed, and in some cases concerns very narrow energy intervals $\Delta E_x \simeq 0.5$ MeV. In this case the use of $\Delta = 0.5$ MeV is unjustified.

2. THE RESULTS OF CALCULATION

The strength functions $C^2(\eta)$, describing the distribution of strength of the deep hole states $1g_{9/2}$, $2p_{1/2}$, $2p_{3/2}$, $1f_{5/2}$ of the ^{111}Sn isotope and of the states $2p_{1/2}$, $2p_{3/2}$, $1f_{5/2}$ of the ^{115}Sn isotope as a function of the excitation energy $E_x = \eta - \eta_{g.s.}$ * are shown in figs. 1 and 2, respectively. Before comparing them with the experimental data, we shall discuss the general assumptions. Table 1 presents certain integral characteristics of distributions

i) the sum spectroscopic factor

$$C^2S = (2J+1)\, v_J^2 \int_0^{E_{max}} C^2(\eta)\, d\eta$$

* $\eta_{g.s.}$ is the ground state energy of an odd-A nucleus calculated with the wave function (1).

ii) the value of the centroid

$$\bar{E}_x = \int_0^{E_{max}} C^2(\eta)\,\eta\, d\eta \Big/ \int_0^{E_{max}} C^2(\eta)\, d\eta - \eta_{g.s.} .$$

iii) the second moment of distribution σ and the so-called full width on half maximum FWHM=2.35σ

$$\sigma^2 = \int_0^{E_{max}} C^2(\eta)(\eta - \bar{E}_x + \eta_{g.s.})^2\, d\eta \Big/ \int_0^{E_{max}} C(\eta)\, d\eta$$

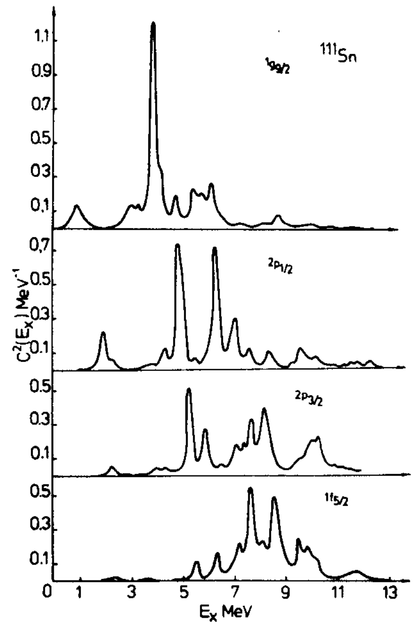

Fig.1. Strength function of the neutron hole states $1g_{9/2}$, $2p_{1/2}$, $2p_{3/2}$ and $1f_{5/2}$ of ^{111}Sn.

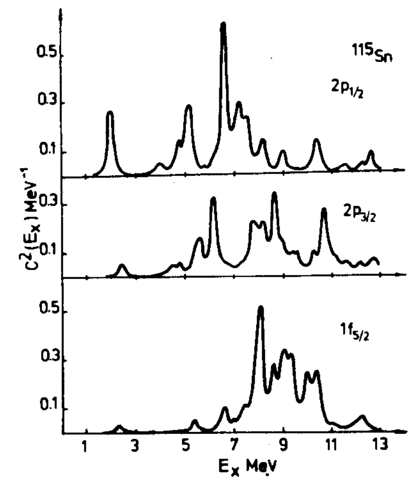

Fig.2. Strength functions of the neutron hole states $2p_{1/2}$, $2p_{3/2}$ and $1f_{5/2}$ of ^{115}Sn.

Table 1 presents also the values of the one-quasiparticle excitation energies of the hole states $\epsilon^*_J = \epsilon_J - \epsilon_{g.s.}$, where $\epsilon_{g.s.}$ is the one-quasiparticle energy of the ground state. The figures and tables show that on the whole the states under consideration are fragmented strongly. The value of FWHM varies within 4-6 MeV (we shall not discuss as yet the state $1f_{7/2}$ of ^{115}Sn). The most strongly fragmented state is $2p_{1/2}$, for which σ is maximal. It is somewhat strange that the state $1f_{5/2}$ has the minimal value of σ, since it has the largest

Table 1

Integral characteristics of strength distribution of deep hole states in ^{111,115}Sn

Nucleus	nlJ	$0\div E_{max}$ MeV	ε_J^* MeV	$\bar{E}_x$ MeV	σ MeV	FWHM MeV	$\sum C^2S$
^{111}Sn	$1g_{9/2}$	0÷13	4.7	4.9	2.0	4.7	9.60
	$2p_{1/2}$	0÷14	6.2	6.3	2.4	5.6	1.96
	$2p_{3/2}$	0÷13	7.6	7.5	2.1	4.9	3.94
	$1f_{5/2}$	0÷14	8.1	8.3	1.6	3.8	5.88
^{115}Sn	$2p_{1/2}$	0÷14	6.7	7.4	2.6	6.1	1.94
	$2p_{3/2}$	0÷14	8.1	8.7	2.3	5.5	3.80
	$1f_{5/2}$	0÷14	8.6	9.3	1.7	3.9	5.88
	$1f_{7/2}$*	0÷20	12.2	13.0	3.0	7.05	7.92

* The strength function for this state has been calculated taking into account the interaction of the hole state with the state of type "quasiparticle plus phonon" alone.

excitation energy than other states. One of the possible reasons of a considerable difference between the values of σ for 2p- and $1f_{5/2}$ -states is a different gross-structure of their strength functions. Almost the whole $1f_{5/2}$ -state strength is concentrated in one excitation energy region with width of about 3 MeV, whereas the strength of other states is concentrated in 2-3 distant intervals of E_x (this is clearly seen especially for the state $2p_{3/2}$). It should be noted that almost the whole strength of the states $1g_{9/2}$, $2p_{1/2}$, $2p_{3/2}$ and $1f_{5/2}$ is exhausted up to the excitation energy $E_x \leq 14$ MeV. The values of centroids $\bar{E}_x$ differ from the corresponding values ϵ_J^* mainly because of the lowering of the ground state energy due to the quasiparticle-phonon interaction. Since in ^{115}Sn the difference of $\epsilon_{g.s.} - \eta_{g.s.}$ is larger than in ^{111}Sn, the difference between ϵ_J^* and $\bar{E}_x$ is also larger in this nucleus. The values of $\epsilon_J^* + \epsilon_{g.s.}$ and $\bar{E}_x + \eta_{g.s.}$ differ by 100-300 keV. The comparison of the values of σ (or FWHM) for the same states of the ^{111}Sn and ^{115}Sn isotopes shows that in ^{111}Sn the fragmentation is somewhat slighter than in ^{115}Sn; this fact has been indicated by experimenters for the state $1g_{9/2}$. One of the reasons of this fact is a larger energy of the hole excitations in ^{115}Sn in comparison with ^{111}Sn.

Table 2

Experimental and theoretical values of the spectroscopic factor C^2S for different excitation energy intervals

Excitation energy intervals (MeV)	C^2S					
	$1g_{9/2}$		$2p_{1/2}$		$2p_{3/2}$	
	exp.	theor.	exp.	theor.	exp.	theor.
3.4÷3.94	1.54	2.97	0.12	0.02	(0.02)	0.03
3.94÷4.5	2.42	1.19	0.12	0.08	0.07	0.08
4.5÷5.85	2.12	1.90	0.22	0.49	0.53	0.74
5.85÷6.80	0.24	1.15	(0.48)	0.47	(0.39)	0.32
6.80÷7.30	0.10	0.13	-	0.21	0.17	0.24
3.4÷5.4	5.37	5.27	$[C^2S(2p_{1/2}) + C^2S(2p_{3/2})]_{exp} = 0.84$ $[C^2S(2p_{1/2}) + C^2S(2p_{3/2})]_{th} = 1.06$			
3.4÷7.3	6.42	7.34	0.46÷ 0.94	1.27	0.8÷ 1.18	1.41

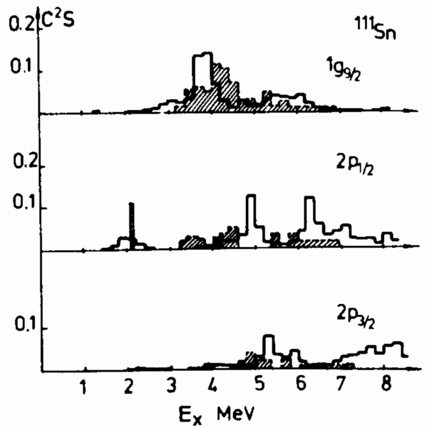

Fig.3. Experimental [5] (shaded) and theoretical histograms of strength distribution of the neutron states $1g_{9/2}$, $2p_{1/2}$ and $2p_{3/2}$ of ^{111}Sn.

Now we proceed with a detailed comparison with experiment. The detailed data for ^{111}Sn pertain to the excitation energy interval 0÷7.3 MeV [5],which is quite sufficient to judge about the fragmentation of the $1g_{9/2}$ subshell and to a great extent of $2p_{1/2}$. This interval includes only the lower one third of the region which comprises the main $2p_{3/2}$ state strength. In paper [5] alongside with the data, which we cite in table 2, the histograms are constructed of the experimental distributions of the quantity C^2S with step 0.2 MeV. In fig.3 we present these histograms together with the theoretical ones constructed according to the data on the strength function $C^2(\eta)$. It is seen from the figures and table that the best agreement between theory and experiment is achieved for the $1g_{9/2}$ state. The difference between the theoretical and experimental results is caused by a lower position (0.2 MeV) of the calculated maximum in comparison with the experimental one and by a higher concentration of the calculated $1g_{9/2}$ state strength in the interval 5.5÷6.5 MeV. The calculated distribution at maximum is narrower than the experimental one. This is confirmed also by the values of the FWHM for the interval 3.4-4.5 MeV; the experimental value of which is 0.575 MeV* and the theoretical one is 0.486 MeV. For the states $2p_{1/2}$ and $2p_{3/2}$, one can also state a qualitative and, to some extent, a quantitative agreement of the theoretical and experimental data. A specific feature of the theoretical distribution of the $2p_{1/2}$ -state strength is the presence of two maxima or the regions of higher concentration of strength: $4.6 < E_x < 5.2$ MeV and $6.0 \lesssim E_x < 6.6$ MeV. Between them there is local maximum of distribution of the $2p_{3/2}$ -state strength (fig.3c). Such a pattern of the 2p -strength distribution is seen in the experimental histogram but with less sharp maxima, which concentrate a smaller value of the single-particle strength and lie by 0.5 MeV below the theoretical ones. The distribution of 2p-strength over narrow bins in comparison with the experimental data is shown in table 2. Treating the total strength of hole states, extracted in the experiments with polarized deuterons, it turns out that the theoretical values of C^2S in the interval 0-7.3 MeV exceed the experimental ones. For the state $1g_{9/2}$ the difference between C^2S_{th} and C^2S_{exp} is 15%. For the subshells $2p_{1/2}$ and $2p_{3/2}$ the experimental errors are large. Comparing C^2S_{th} with the upper bound of the experimental value of the spectroscopic factor, for the state $2p_{1/2}$ the difference is 25%; and for $2p_{3/2}$, 16%.

To separate in the cross sections of the one-nucleon transfer reactions the contribution of deeper hole states is not a simple task. The quantitative results obtained here depend considerably on the subtracting background. The reason is that with increasing excitation energy, the fragmentation of single-particle states becomes stronger, the widths of the corresponding structures increase and they are overlapped strongly.

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*The experimental value of FWHM is taken from ref. /15/. In ref. /5/ the value of σ is not presented, and the difference of the results of refs. /5/ and /15/ is not large.

A well localized resonance-like peak in the experimentally measured spectrum exists just for the excitation of the state $1g_{9/2}$; the 2p -shell strength is already fragmented strongly. Greater difficulties and uncertainties should arise when passing to the state 1f. The data on the excitation of the subshell $1f_{5/2}$ have been obtained in paper /4/, in which the excitation energy region of ^{115}Sn investigated in the reaction (d, t) (E_d=50 MeV) has been extended up to E_x =25 MeV. This allowed one to determine anew the background by diminishing it by 40% in comparison with earlier measurements, in which a narrower region E_x has been investigated /16/. Owing to this the tails of distributions of the hole states $1g_{9/2}$, 2p and $1f_{5/2}$ have been distinguished. It was shown that in the interval $E_x \le 20$ MeV the whole strength of these states is exhausted. The corresponding data for the 2p- and 1f-states are presented in table 3. This table also presents the theoretical sum values C^2S for the same excitation energy intervals*. A satisfactory agreement with experiment is obtained for the 2p -states. There is an agreement with the data of another group /17/. They refer to a narrower interval E_x and are of a tentative nature. It is pointed out that in ^{115}Sn in the interval $4.9 < E_x < 5.9$ MeV an increased concentration of the $2p_{1/2}$ -strength is observed, and $C^2S \simeq 0.3$. It is seen from fig.2 that a local maximum of the strength distribution of this state is located in this region. The value of C^2S_{th}=0.28. For the states 1f the results of our calculations differ from the estimates obtained from the experimental data. According to the data of paper /4/ the centroid of the distribution for the state $1f_{7/2}$ is at an energy of 10.6 MeV, and FWHM ≈ 8 MeV. These values do not coincide with our data presented in table 1. A considerably less theoretical value of FWHM may partly be due to the restrictions of the calculations, which have been mentioned in this paper. A considerable difference $\bar{E}_x(\text{exp.}) - \bar{E}_x(\text{th.}) = 1.3$ MeV is somewhat surprising. The position of E_x(th) is determined mainly by the single-particle scheme, and our values of the parameters of the Saxon-Woods potential /18/ have been checked in many calculations. Besides, the authors of paper /4/ concider that the whole extracted strength of $\ell = 3$ transitions is due to the $1f_{5/2}$ subshell excitation. The experimental values of C^2S, presented in

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*It should be mentioned that the boundaries of intervals in the table are determined roughly according to fig.1 in ref. /4/, since in the table of the same paper they are not presented.

Table 3

Strength distribution of the 2p- and 1f-subshells in the spectrum of ^{115}Sn

Excitation energy intervals (MeV)		C^2S Exp.* $\ell=1$	Theor. $2p_{1/2}$	Theor. $2p_{3/2}$	Theor. $\sum 2p$	Exp.* $\ell=3$	Theor. $1f_{5/2}$	Theor. $1f_{7/2}$**	Theor. $\sum 1f$
I+II	0÷5.5	0.59±0.04	0.35	0.27	0.62	-	0.15	0.09	0.24
III	5.5÷7.5	2.0±0.2	0.70	0.90	1.60	-	0.40	0.10	0.50
IV	7.5÷10.5	1.2±0.3	0.62	1.67	2.29	-	3.9	1.16	5.06
V	10.5÷19.5	1.9±0.4	0.3	1.15	1.45	8.6±1.4	1.5	6.25	7.75

*For ℓ=1 transitions, experimental C^2S -values obtained for $2p_{1/2}$ and $2p_{3/2}$ pickup were averaged. For ℓ = 3, $1f_{5/2}$ pickup was assumed by experimenters.

**Calculation is performed taking into account the interaction of the hole state with the "quasiparticle plus phonon" states alone.

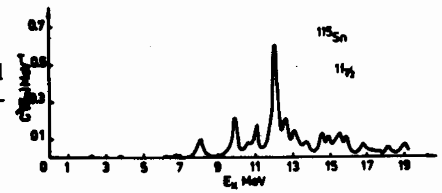

Fig.4. Strength function of the state $1f_{7/2}$ in ^{115}Sn. The calculation has been performed taking into account the interaction of the state $1f_{7/2}$ with the "quasiparticle plus phonon" states alone.

table 3, are obtained just under this assumption. This is not verified in our calculations. Figure 4 shows the strength function for the state $1f_{7/2}$ of ^{115}Sn. Because of great computational difficulties, arising in the calculations with the wave function (1) at so large excitation energies, we have used in (1) the components of type "quasiparticle plus phonon". The data specifying the strength distribution of the $1f_{7/2}$ state are given in tables 1 and 3. The distribution centroid is at the energy $\bar{E}_x$ =13 MeV, and the $1f_{7/2}$ state should give a noticeable contribution to the cross section of the one-nucleon transfer reaction, starting from the excitation energy $E_x \simeq$ 10 MeV. The inclusion of the "quasiparticle plus two phonons" components, enhancing the fragmentation, may lower this boundary. It should be noted that the state $1f_{7/2}$ is very strongly fragmented. Our calculations give, apparently, the lower boundary of the fragmentation of this state.

CONCLUSION

Based on the above-mentioned results of calculation we can make the following conclusions: the quasiparticle-phonon nuclear model describes satisfactorily the available experimental data on the fragmentation of deep-lying hole states in the tin isotopes. For the excitation energies $E_x \leq 8$ MeV there is a fairly good quantitative agreement of theory with experiment, though in the calculations the fragmentation is systematically weaker than in the experiment. At higher excitation energies $E_x > 10$-15 MeV the theoretical results and experimental data differ more strongly. However, we think that it is untimely to make pretension to the model. In this region of energies the simple one-hole states are so strongly fragmented, that it is very difficult to evaluate reliably their strength by the data of the one-nucleon transfer reaction.

The authors are deeply indebted to Drs. S.Gales, E.Gerlic and H.Langevin-Joliot for communicating to us their results prior to publication and for useful discussions.

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Received by Publishing Department
on April 5 1982.