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STRENGTH FUNCTIONS OF HIGH-LYING SINGLE-NEUTRON STATES

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Since the initial observations of deep-hole states in neutron pick-up reactions in the early 70th [1] the one-nucleon transfer reaction has been widely used to study the high-lying single-particle (or single-hole) strengths in medium and heavy nuclei. Very rapidly a growing amount of experimental data has been accumulated and new exciting features of nuclear excitation spectra have been found. In the proton pick-up reactions the deep-hole proton states have been observed [2] and in the proton stripping reactions the high-lying single-proton states have been found [3]. Now, the high energy resolution gives a possibility of studing the fine structure of excited resonance-like structures, and using a polarized incident beams experimenters can make unambiguous spin assignments to it.

So, the one-nucleon transfer reactions at high energies made it possible to investigate the new kind of nuclear excitations which one may call "giant single-particle (or single-hole) resonances". These giant resonances are the response of a nucleus on the external field acting on its single-particle degrees. The comparison of the experimental data with the predictions of different theoretical models helps us to check our knowledge on the average nuclear field and on the mechanisms responsible for the spreading of the single-particle strength in nuclei. On the other hand, to extract the single-particle (or hole) strength from an experimental cross-section it is useful to have some theoretical predictions about the fragmentation of high-lying nuclear subshells because their strength distributions are overlapped.

One of the last experimental achievements is the observation of high-lying single neutron states in tin isotopes with mass-numbers $A = 117$, 119 and 121 [4]. Analogous experiments for the target-nucleus ^{90}Zr, ^{144}Sm, and ^{208}Pb are now in progress [5]. So, we try to calculate the strength functions of these high-lying single-neutron subshells in the corresponding N-odd nuclei. For the tin isotopes we compare our results with the available experimental data.

We use the quasiparticle-phonon nuclear model and the strength function method in our calculations. The quasiparticle-phonon nuclear model (QPM) has been expounded in detail in a series of reviews [6-9]. The residual NN-interaction in the QPM consists of the traditional monopole pairing interactions in the particle-particle channel and separable multipole and spin-multipole interaction in the particle-hole channel. The particle-hole interaction is isotopically invariant. It is well known now that damping widths of giant resonances and resonance-like structures in nuclei are due to the coupling of simple nuclear modes (in our case - single-particle modes) with the complex configurations. The coupling with low-lying surface vibrations plays the key role in the spreading of the strengths of simple excitations [8-11]. But our previous studies [9,10] have shown the importance of the coupling with collective phonon excitations at intermediate energies (e.g., M2-resonance or LEOR) to describe the fine structure of the strength distributions. In the framework of the QPM we take into account the interaction of single-particle states with all types of phonons: low-lying vibrations, collective phonons of intermediate energies and noncollective phonons (i.e., almost pure two-quasiparticle states). We should like to stress that the quasiparticle-phonon interaction H_{qph} is calculated microscopically in the QPM. The interaction H_{qph} depends on the characteristics of phonons which are calculated in the RPA and on single-particle matrix elements of residual particle-hole interaction. In terms of the creation and annihilation operators of quasiparticles and phonons (a^+_{jm}, a_{jm} and $Q^+_{\lambda\mu i}$, $Q_{\lambda\mu i}$) the operator H_{qph} has the form:

$$H_{qph} = -\frac{1}{2\sqrt{2}} \sum_{\lambda\mu i} \{(Q^+_{\lambda\mu i} (-)^{\lambda-\mu} + Q_{\lambda-\mu i}) \sum_{jj'\tau} \frac{f^{(\lambda)}_{jj'} v^{(-)}_{jj'}}{\sqrt{\mathcal{Y}^{\lambda i}_{\tau}}} B(jj'\lambda-\mu) + h.c.\}, \tag{1}$$

$$B(jj'\lambda-\mu) = \sum_{mm'} (-)^{j'+m'} \langle jmj'm'|\lambda-\mu\rangle a^+_{jm} a_{j'-m'}, \tag{2}$$

where $f^{(\lambda)}_{j_1 j_2}$ is reduced single-particle matrix element of the operator $f(r)Y_{\lambda\mu}$ *; $v^{(-)}_{j'j} = u_{j'}u_j - v_{j'}v_j$ (u_j, v_j - the coefficients of Bogolubov's transformation). The factor $\mathcal{Y}^{\lambda i}_{\tau}$ is normalizing factor in the one-phonon wave function. There are two of them for each phonon: one for the neutron part of the wave function ($\tau = n$) and the other for the proton one ($\tau = p$). The factors $\mathcal{Y}^{\lambda i}_{\tau}$ depend on the collectivity of the phonon λi. The values of $\mathcal{Y}^{\lambda i}_{\tau}$ are small for the collective phonons and large for the noncollective ones. Note, the same isotopic index τ as in $\mathcal{Y}^{\lambda i}_{\tau}$ is included in the single-particle quantum

* In this work $f(r) = dU/dr$ and U denotes a central part of the Saxon-Woods potential.

number $j \equiv n\ell j\tau$. One can see from (1) that H_{qph} couples the states which differ from each other by one phonon, i.e.,

$$a^+_{JM}\Psi_0 \to [a^+_{jm} Q^+_{\lambda\mu i}]_{JM}\Psi_0 ;$$

$$[a^+_{jm} Q^+_{\lambda\mu i}]_{JM}\Psi_0 \leftrightarrow [a^+_{j'm'} [Q^+_{\lambda_1\mu_1 i_1} Q^+_{\lambda_2\mu_2 i_2}]_{IM_1}]_{JM}\Psi_0 .$$

When calculating the matrix element of H_{qph} we don't take into account the anharmonic corrections for the phonon excitations. We believe these corrections affect slightly the fragmentation of high-lying single-particle strength. Moreover we use a special approximate procedure to preserve the violation of the Pauli principle in the complex components of our model wave function. We have studied all these effects correctly in refs. [9,12,13]. But to solve numerically the full set of equations of the QPM [13] for high-lying states in odd-mass nuclei is a heavy computational problem. At present we use the model wave function which consists of one-quasiparticle, "quasiparticle ⊗ one-phonon" and "quasiparticle ⊗ two phonons" components and solve simpler set of equations than in ref. [13] (see, ref. [12]). We take into account the coupling of an odd-quasiparticle with all the phonons with momenta and parities $\lambda^\pi = 1^\pm \div 7^\pm$ and excitation energies less than 20 MeV.

We don't calculate the energy and wave function structure of each of the numerous high-lying excited states. Instead we calculate the single-particle strength function $C^2_j(E_x)$ which describes the dependence of the smearing on the energy interval Δ the square of one-quasiparticle amplitude C^2_j on the excitation energy E_x. The strength function method has been used firstly by Bohr and Mottelson [14] for a schematic model. Later the method has been developed for more complex problems [15] (see also ref. [8,9]) and now it is widely used in nuclear structure calculations [11]. The strength function method makes the calculation much simpler and the results more obvious. In the present work to construct a single-neutron strength function we use as a weight function the Lorentz function with the width parameter Δ = 0.5 MeV.

We have discussed the model parameters in detail in refs. [7,8,10]. Here we display in fig.1 only the parts of single-neutron schemes of the nuclei in question. The single-particle energies and wave functions for all the nuclei except ^{209}Pb are calculated in the spherically-symmetric Saxon-Woods potential with standard parameters from refs. [16]. These parameters have been used in our previous calculations. The energies of the levels of the shells which are closest to the Fermi surface in ^{209}Pb are taken from ref. [17]. In that paper they have

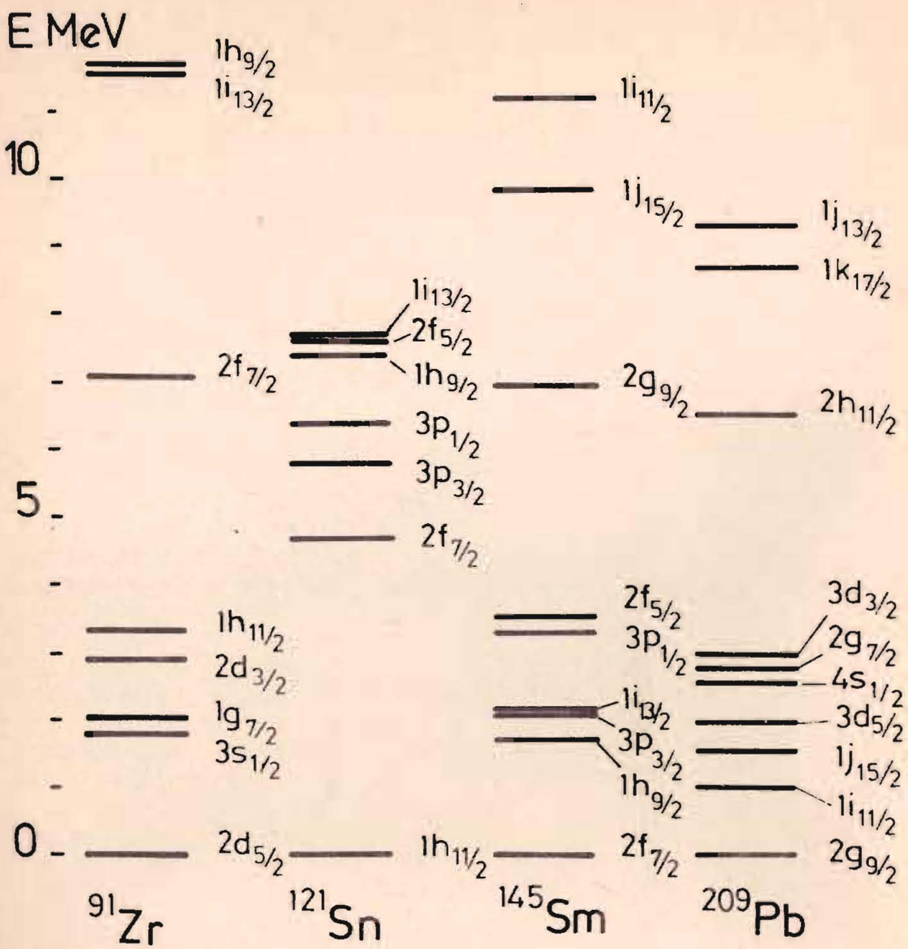


Fig. 1. Parts of single-neutron schemes in nuclei ^{91}Zr, ^{121}Sn , ^{145}Sm, *and* ^{209}Pb. *The Fermi level is assumed to have zero energy.*

been fitted in the QPM calculation to the experimental data on the energies and one-nucleon spectroscopic factors of low-lying states of the nuclei ^{207,209}Pb, ^{207}Tl, and ^{209}Bi. The energies of single-particle (and hole) states well above and below the Fermi surface are calculated in the Saxon-Woods well too.

We should like to start the discussion of our results with the tin isotopes. The clearly seen resonance-like structures at excitation energies 4-10 MeV have been observed in the study of the ^{116,118,120}Sn(α,^{3}He) reactions at 183 MeV incident energy /4/. The measured excitation energies and angular distributions sup-

Experimental data and theoretical results for the strength distributions of $1i_{13/2}$ *and* $1h_{9/2}$ *-subshells in* ^{121}Sn

	Experiment /4/			Theory		
	E_m, ΔE_x(MeV)	J^π	C^2S	E_m,ΔE_x(MeV)	J^π	C^2S
<A>	E_m =4.89±0,05	9/2⁻	0.67	E_m = 5.1 ΔE_x = 4,6-5,6	9/2⁻	0,12
		–	–		13/2⁺	0,21
<B>	ΔE_x=5.4-7.0	9/2⁻	0,25	ΔE_x=5.4-7.0	9/2⁻	0,4
		13/2⁺	0,16		13/2⁺	0,20
<C>	ΔE_x=7,0-10,0	–	–	ΔE_x=7.0-10.0	9/2⁻	0,34
		13/2⁺	0.27		13/2⁺	0,39

E_m – *the energy of the peak;* ΔE_x – *the energy interval for which the value* C^2S *is founded.*

port the assumption that these bumps arise from neutron stripping to the $1h_{9/2}$ and $1i_{13/2}$ orbitals located well above the valence subshell. When the mass number of the isotope decreases, the centroid energy and the spreading of the strength distribution increases. In ^{121}Sn, the resonance-like structure consists of a sharp peak at E_x = 4.89 MeV with a width less than 1 MeV (a range <A> in table 1) and additional broader components (<B> and <C> in table 1). In ^{119,117}Sn this narrow peak smears and almost disappears. The quantitative data on the single-neutron strength distributions are only for ^{121}Sn.

In the one-nucleon transfer reaction at intermediate incident energies the transitions with large transfer momenta are enhanced. Therefore, at first the strengths of the subshells with high orbital momentum are extracted from experimental cross-sections. In tin isotopes those neutron subshells are $1h_{9/2}$ and $1i_{13/2}$. These quasibound neutron states are located in the single-neutron spectrum very closely to each other. As one can see in fig.2 and table 2 in our calculation the main strengths of both the subshells are spread over the same energy interval $4 \leq E_x \leq 9$ MeV. The fragmentation of the subshells in

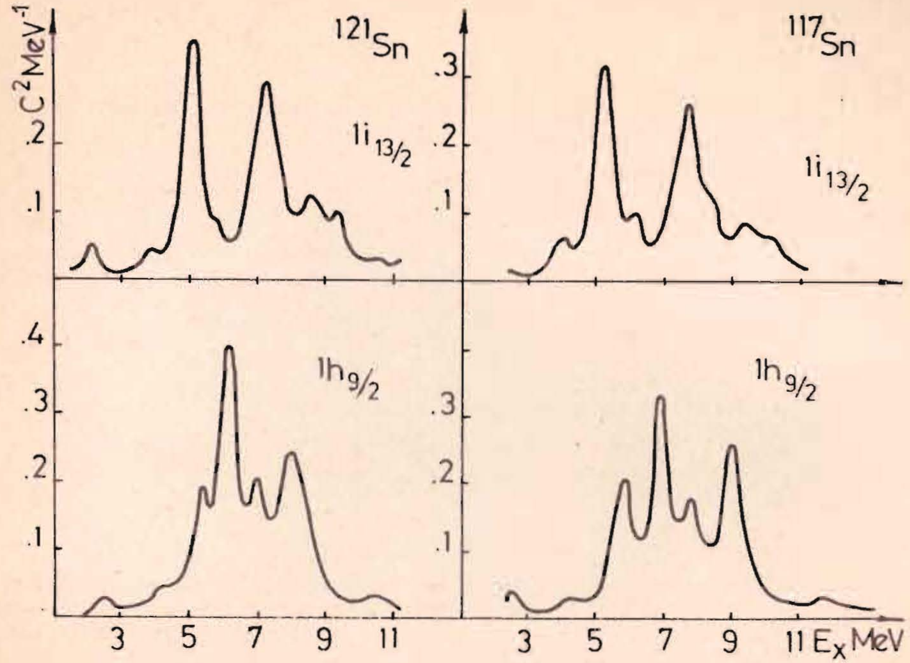


Fig.2. Strength functions of the high-lying neutron states $1i_{13/2}$ *and* $1h_{9/2}$ *in* ^{117,121}Sn.

^{117}Sn is stronger than in ^{121}Sn to some extent. This difference is the most prominent for the subshell $1h_{9/2}$. The strength function maxima are located at $E_x \sim 5 \div 6$ MeV therefore the relatively high subshell strengths would be observed here. So, it seems to us that there is a qualitative agreement between the experimental data and our results.

But the quantitative agreement is somewhat worse (see table 1). The most striking, although not so important, disagreement is for the energy interval <A>. According to the experimental interpretation, at this region almost 2/3 of the total $1h_{9/2}$-strength is concentrated. The maximum of the theoretical $1h_{9/2}$-strength function is placed at E_x = 6.2 MeV, i.e., by 1.3 MeV higher than the experimental peak. But there is main peak of the $1i_{13/2}$-strength function (E_x = 5.1 MeV) very close to the experimental one. Of course, the accuracy of the theoretical model is not high enough to make a definite conclusion on the fine structure of the strength distributions of high-lying states. However, some part of the $1i_{13/2}$-strength must be located in the region <A> of the spectrum. We should like to point out that at scattering angles $\theta > 10°$ the experimental angular distribution is closer to the theoretical one calculated

Table 2

The theoretical characteristics of the high-lying single-neutron strength distributions in spherical nuclei ($\bar{E}$ *is the energy centroid,* σ *is the second moment,* C^2S *is the part of the total state strength which is exhausted in the interval* ΔE_x *)*

Nucleus		ΔE_x, MeV	$\bar{E}$, MeV	σ, MeV	C^2S
^{91}Zr	$2f_{7/2}$	2.5 - 13.5	8.1	1.76	0.93
		7.2 - 9.7	8.45	0.75	0.55
	$1i_{13/2}$	5.5 - 18.5	12.1	2.19	0.85
		9.4 - 13.9	11.5	1.22	0.54
	$1h_{9/2}$	7.5 - 16.5	12.6	1.60	0.83
		10.7 - 13.5	11.9	1.41	0.51
		13.5 - 16.5	14.9	0.95	0.24
^{117}Sn	$1i_{13/2}$	2.0 - 11.0	6.9	1.9	0.79
		4.3 - 6.8	5.4	0.53	0.30
		6.8 - 8.9	7.8	0.53	0.32
	$1h_{9/2}$	2.0 - 13.0	7.5	2.0	0.91
^{121}Sn	$1i_{13/2}$	1.0 - 11.0	6.7	2.0	0.89
		3.0 - 6.0	5.0	0.57	0.31
		6.0 - 8.2	7.3	0.48	0.33
	$1h_{9/2}$	1.0 - 11.0	6.8	1.6	0.92
^{145}Sm	$1i_{11/2}$	5.0 - 16.0	12.3	1.94	0.90
		10.0 - 12.6	11.5	0.46	0.45
	$1j_{15/2}$	5.0 - 16.0	10.9	2.46	0.87
		7.0 - 10.0	8.7	0.72	0.31
		10.0 - 12.0	11.1	0.46	0.23
^{209}Pb	$1k_{17/2}$	2.5 - 15.5	8.9	2.32	0.94
	$1j_{13/2}$	2.5 - 15.5	9.6	1.8	0.94

for transfers ℓ = 6, although at $\theta < 10°$ the transfers ℓ = 5 are preferable.

A more important disagreement concerns the wide energy intervals ΔE_x. From data of ref. /4/ (table 1) it is seen that almost the total $1h_{9/2}$-strength (92%) is exhausted in the interval ΔE_x ~2 MeV ($4.9 \le E_x \le 7$ MeV). At the same time only

43% of the $1i_{13/2}$-strength has been observed at wider interval $5.4 \leq E_x \leq 10$ MeV. From these data one can make a conclusion that the $1i_{13/2}$-subshell has the excitation energy much higher than that of the $1h_{9/2}$-subshell and is fragmented stronger. These conclusions contradict our results. Of course, the positions of the $1h_{9/2}$ and $1i_{13/2}$-single-neutron states have been calculated in the phenomenological potential well and one can try to change potential parameters to improve the agreement with new experimental data. But the close positions of these subshells in a single-neutron scheme is a rule for all the nuclei in question (see fig.1). Moreover, so strong concentration of $1h_{9/2}$-strength in the excitation spectrum seems to us unlikely.

In view of these results we should like to stress once more that to extract the subshell strengths from experimental cross-sections correctly it is necessary to take into account the overlapping of their distributions. In some cases the contribution of subshells with relatively low orbital momentum numbers becomes important. For example, the agreement of theoretical and experimental strength distributions of the $1h_{9/2}$- and $1i_{13/2}$-single-proton states had been improved noticeably since the contribution of the $2f_{7/2}$-subshell to the one-proton stripping cross-section was taken into account [3,18]. It means that in the present case experimenters would take into account the $2f_{5/2}$- and $2f_{7/2}$-strengths. The calculations of the strength functions of these subshells are in progress.

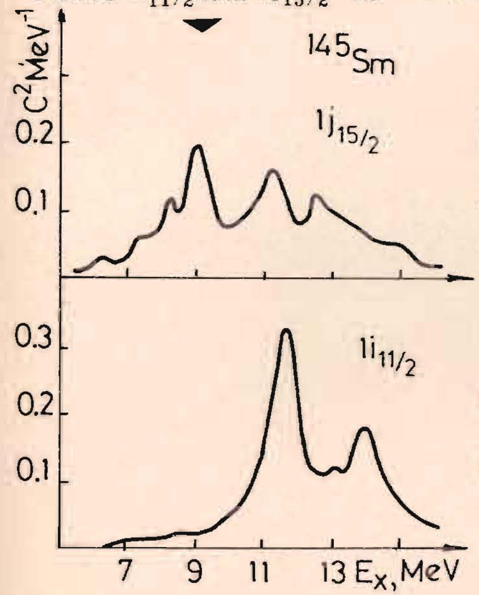
The calculated strength functions of the high-lying single-neutron states in ^{91}Zr, ^{145}Sm and ^{209}Pb are displayed in figs. 3,4,5. The integral characteristics of the single-neutron strength distributions are given in table 2. Now only preliminary experimental data on these states are available [5] and we hope that our calculations will help to analyze the one-neutron stripping reaction cross section. The results for ^{91}Zr have been pulblished in ref. [19].

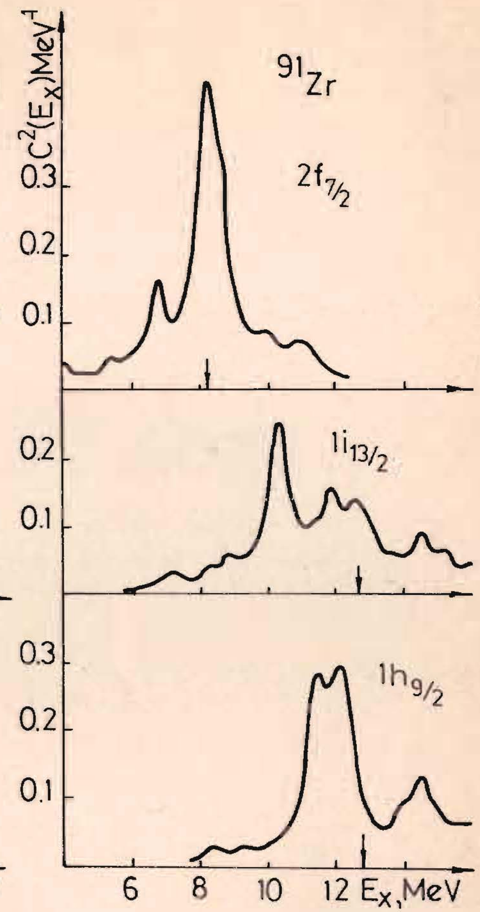
In all three nuclei - ^{91}Zr, ^{145}Sm and ^{209}Pb - an odd neutron is coupled with even-even core with magic neutron number (N = 50, 82, and 126, respectively). The valence shells begin to occupy and the excitation energies of quasibound neutron states is higher than in $^{117\text{-}121}$Sn. Therefore, the escape widths of the quasibound states with low orbital momentum is large and these states are not shown in figs.3,4,5. The displayed subshells have small escape widths. The patterns of the quasibound part of the single-neutron scheme in ^{91}Zr, ^{145}Sm and ^{209}Pb are similar. Two subshells with high ℓ-numbers* are located close to each other and there is a subshell with much less ℓ -

* In ^{91}Zr $\ell = 6$; in ^{145}Sm $\ell = 7$ and in ^{209}Pb $\ell = 8$.

Fig.3. Strength functions of the high-lying neutron subshells $2f_{7/2}$, $1i_{13/2}$ *and* $1h_{9/2}$ *in* ^{91}Zr. *Arrows point out the energies of the one-quasiparticle states* $2f_{7/2}$, $1i_{13/2}$ *and* $1h_{9/2}$.

Fig.4. Strength functions of the high-lying neutron subshells $1i_{11/2}$ *and* $1j_{15/2}$ *in* ^{145}Sm.





number by several MeV below them. Note, that there are the narrow quasibound neutron states with higher orbital momentum in single-neutron spectra, but they are located at much higher excitation energy. We believe that due to very strong spreading the corresponding resonance-like structures will be hardly observable.

Our calculations show that all the subshells studied are fragmented strongly. The second moments of the strength distributions are in the interval $1.6 \leq \sigma \leq 2.5$ MeV. The strength distributions of the high-ℓ subshells are overlapped and this effect should be taken into account in analyzing experimental data. The effect of the low-ℓ subshells will be weaker than in tin isotopes because they are located well below the high-ℓ quasibound subshells. We assume that it is possible to observe

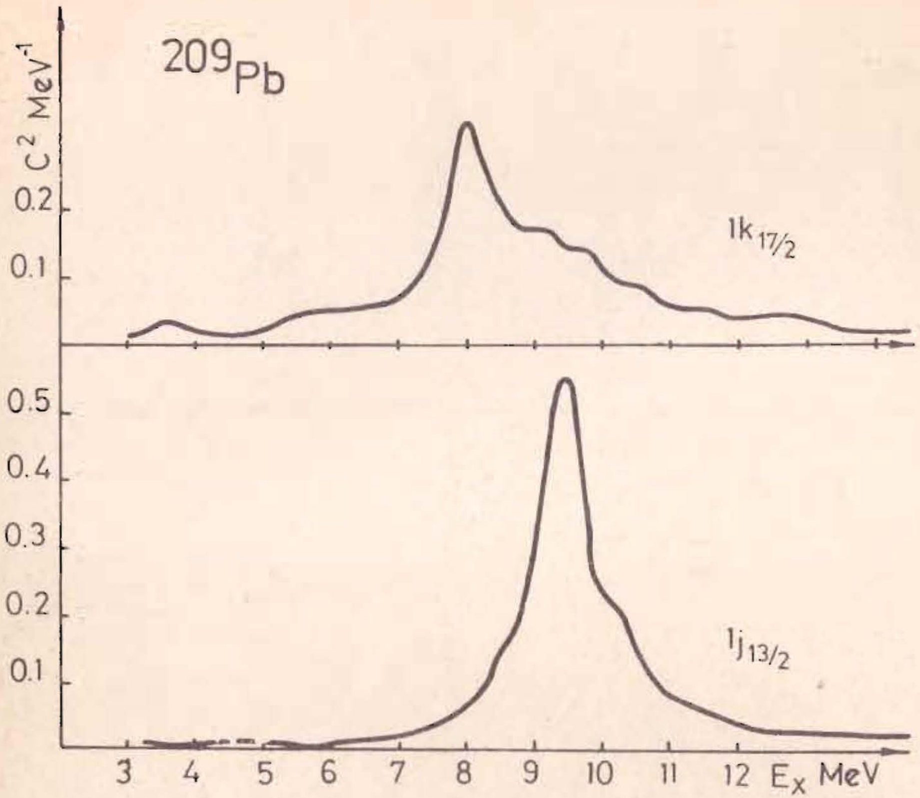


Fig.5. Strength functions of the high-lying neutron subshells $1j_{13/2}$ *and* $1k_{17/2}$ *in* ^{209}Pb.

the resonance-like structures which are due to the one-neutron stripping on the subshells $2f_{7/2}$, $2g_{9/2}$ and $2h_{11/2}$ in ^{91}Zr, ^{145}Sm, and ^{209}Pb, respectively.

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Вдовин А.И., Стоянов Ч. E4-85-352
Силовые функции высоколежащих нейтронных одночастичных состояний

В рамках квазичастично-фононной модели ядра рассчитаны силовые функции одночастичных нейтронных состояний с энергиями возбуждения 7-12 МэВ в ядрах ^{91}Zr , ^{117,121}Sn , ^{145}Sm и ^{209}Pb. Результаты расчетов для изотопов олова сравниваются с экспериментальными данными.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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Vdovin A.I., Stoyanov Ch. E4-85-352
Strength Functions of High-Lying Single-Neutron States

Strength functions of single-neutron states with excitation energy from 7 to 12 MeV in nuclei ^{91}Zr , ^{117,121}Sn , ^{145}Sm , and ^{209}Pb are calculated. The quasiparticle-phonon nuclear model and strength function method are used. The results of calculations are compared with available experimental data for tin isotopes.

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

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