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QUASIFREE A(p, p′d) B LARGE ANGLE SCATTERING AT INTERMEDIATE ENERGIES

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INTRODUCTION

The quasifree knockout reactions have a distinguished role in the investigation of clustering aspects of nuclear structure. The incident particle can be rescattered by a cluster in the nucleus and can give the cluster as a whole a momentum sufficient for knocking it out of the nucleus. The high probabilities of the quasielastic knockout of deuterons or alphas have suggested the presence of a large number of preformed clusters in nuclei.

According to this picture it is generally accepted that only those nucleons of the target A contribute significantly to the cross section of the A(a,ab)B quasifree scattering whose relative motion and spin-isospin wave function is identical to those of the emitted cluster b [1]. Calculating an overlap integral we project the clusters b from the target A and in the final formula of the A(a,ab)B cross section there appears the product of a free $ab \to ab$ scattering cross section and the effective number of clusters in the target A .

Since the first step of the development of the theory it has been emphasized that the "noncluster" components of the target wave function, for example in exchange processes, can give large contributions to the quasifree processes [2]. The knockout processes due to the rearrangment of the virtual cluster in the nucleus in the course of its knockout may have an important role, too [3].

All in all the existence of "preformed" alphas in nuclei seems to be a well established concept [4], but there are objections against the idea of "preformed" deuterons in nuclei [5]. One of the serious objections can be obtained from the evaluation of the quasifree knockout reactions on nucleus ^{6}Li. The distorted wave impulse approximation (DWIA) analysis of different ^{6}Li(a,a'd)^{4}He processes using $d-\alpha$ cluster wave function has given different effective numbers of deuteron in the p-shell, in a wide spectrum of values from 0.3 to 1.75 [6,7]. The accepted theoretical value is at 0.5 ~ 0.6 [8]. The explanation has been sought in terms of contraction of deuteron cluster depending on the momentum value P of the residual nucleus ^{4}He [9]. The calculation of the rms value

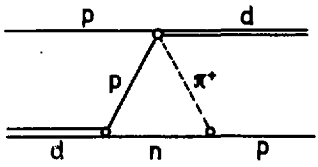

Fig.1. The OPE triangle diagram for the $pd \to pd$ scattering.

of the pn cluster in nucleus ^{6}Li using the three-body wave function[10] shows a drastic shrinkage of the pn cluster radius strongly depending on p[11].

It seems to be an acceptable conclusion that the concept of the "preformed" deuteron cluster is crude in some cases. If the exchange processes are important then the cross section of the $a\langle pn\rangle \to ad$ scattering ($\langle pn\rangle$ is the $p-n$ subsystem inside the nucleus) depends on the internal motion of the $\langle pn\rangle$ cluster only and not on the similiarity with the deuteron defined by an overlap integral of type $\int \phi_d^*(\vec{r})\phi_{\langle pn\rangle}(\vec{r})\, d\vec{r}$. We have to note that the knowlegde of the mechanism of the $a\langle pn\rangle \to ad$ process is generally poor and the formalism of the $A(a,ad)B$ reaction based directly on the $a\langle pn\rangle \to ad$ amplitudes can be very complicated[12].

In the present paper we discuss a model for the description of the $A(p,p'd)B$ quasifree large angle scattering at intermediate energies. For the amplitude of the elementary process $p\langle pn\rangle \to pd$ we have applied the triangle mechanism of[13]. Using the DWIA analysis we can give the final $A(p,p'd)B$ cross section in a relatively simple form.

THE $p\langle pn\rangle \to pd$ AMPLITUDE

As the $p\langle pn\rangle \to pd$ amplitude can not be studied directly on the structure of the $p\langle pn\rangle \to pd$ amplitude we can get information by studying the free $pd \to pd$ scattering. There are different models for the pd large angle scattering at intermediate energies[13-17], which describe well the differential cross sections. All of them contain in some form the Δ_{33} resonance which is created in the NN collisions at incident energie $T_0 \sim 620$ MeV with the width $\Gamma \sim 120$ MeV. The calculated polarizations within the triangle diagram of Craigie and Wilkin (see fig.1) give a good agreement with the experiments for the vector polarization in the 500-700 MeV[18] and for the tensor polarization in the 400-1000 MeV region[19].

In addition, the triangle model is the simplest and so this model is convinient for our further investigation. The pd differential cross section due to the one-pion-exchange (OPE) diagram in fig.1 is proportional to the $pp \to d\pi^+$ cross section which can be taken from the experiment. There are different versions of the triangle model and we have chosen the method of [14], as their work describes the dpn vertex in terms of nonrelativistic deuteron wave functions. This model does not contain free parameters.

The $pd \to pd$ amplitude due to the OPE diagram depends on the initial $\phi_{d\ell}(r)$ deuteron wave function with components $\ell = 0$ or $\ell = 2$ in the following form [14,20] :

$$a_\ell \sim \int_0^\infty \phi_{d\ell}(r)\, e^{-\gamma r}\, (1+\gamma r)\, j_1(\tilde{p} r)\, dr\,, \qquad /1/$$

$$\gamma^2 = T_2^2(1+T_2/m)^{-2} + \mu^2/(1+T_2/m)\,, \qquad /1a/$$

$$\tilde{p} = p_2 (1+T_2/m)^{-1}, \qquad /1b/$$

where m and μ are the nucleon and pion masses, respectively, T_2 is the kinetic energy of the final proton in the laboratory system, $p_2 = (2mT_2 + T_2^2)^{1/2}$; $j_\ell(x)$ is a spherical Bessel function. In the isobar region $\gamma \sim 0.75\ fm^{-1}$ and $\tilde{p} \sim 1.5\ fm^{-1}$. This means that the a_ℓ amplitude depends on the short range part of the deuteron wave function only, the most sensitive region is about 1 fm.

The basic approximation in derivation of a_ℓ is the peaking approximation, that is, the $pp \to d\pi^+$ amplitudes and other slowly variable factors in the $N \to N\pi$ vertex are replaced by their value at zero pn relative momentum in the deuteron. The peaking approximation is quite accurate for the s-wave component but the d-wave contribution is strongly supressed if we have an accurate treatment [19].

Accepting the OPE model for the free pd large angle scattering we can extend it to the $p\langle pn\rangle \to pd$ reaction. In the case of the $\langle np,\, t=0\rangle$ two nucleon cluster (t is the isospin quantum number) the generalization means simply the exchange of the $\phi_{d\ell}(r)$ function in (1) to the $\phi_{\langle np\rangle\ell}(r)$ function. The extension for $\langle pn, t=1, \nu\rangle$ systems (ν is the isospin projection quantum number) can be done considering the isospin invariance [20].

Extending the formalism for higher ℓ orbital momenta we find that the main feature of the a_ℓ amplitudes is the dominance of a_0 . First the sensitivity of the amplitude on the short range part of the wave function prefers the components with lower ℓ due to the repulsing effect of the centrifugal potential [20]. Secondly in the accurate treatment of the OPE diagram the $\ell \neq 0$ components are supressed [19].

The differential cross section of the $p \langle pn, t=0 \rangle \to pd$ reaction can be written with a small change of the $pd \to pd$ cross section given in paper [20].

$$\frac{d\sigma^{p\langle pn\rangle \to pd}}{d\Omega_\Phi} = \frac{3}{2}\,\frac{G^2}{4\pi}\,F^2(k^2)\,\frac{E_2+m}{E_2^2}\,\frac{s_{pp}}{s_{pd}}\,\frac{|p|}{|d|}\,\frac{3}{2}\,\frac{d\sigma^{pp\to d\pi^+}}{d\Omega_\Theta}\,f_0^2 \tag{2}$$

$$= A^2(s_{pd}, u_{pd})\, f_0^2$$

with

$$f_0 = \int_0^\infty \phi_{\langle pn\rangle 0}(r)\, e^{-\gamma r}\,(1+\gamma r)\, j_1(\tilde{p} r)\, dr \tag{2a}$$

$G^2/4\pi = 14.8$. $F(k^2)$ is the Ferrari-Selleri factor [21] which takes into account the off-mass shell nature of the pion. $E_2 = T_2 + m$, $A^2(s_{pd}, u_{pd})$ is defined by the equation (2). The four momentum quadrats s_{pp}, s_{pd}, $|p|$, $|d|$, k^2, u_{pd} can be expressed as follows:

$$s_{pd} = (p_0 + d_0)^2 , \tag{3a}$$

$$s_{pp} = (s_{pd} - m^2)/2 , \tag{3b}$$

$$|p|^2 = \frac{1}{4} s_{pp} - m^2 , \tag{3c}$$

$$|d|^2 = \frac{1}{4 s_{pp}} [(s_{pp} - m_d^2 - \mu^2)^2 - 4 m_d^2 \mu^2] , \tag{3d}$$

$$k^2 = \frac{1}{2} m^2 + m_d^2 - \frac{1}{2} s_{pd} + |p|^2 (1 - \cos\phi) = - m_d T_2 , \tag{3e}$$

$$u_{pd} = m^2 + 2k^2 , \tag{3f}$$

where m_d is the deuteron mass, p_0 and d_0 are the four momenta of the proton and the deuteron in the initial state, respectively. For the sake of simplicity the binding energy of the <pn> system is assumed to be equal to that of the deuteron. The $pp \to d\pi^+$ differential cross section taken at the angle Θ corresponds to the $\cos\Theta$ fixed prescription /15/ the relation between the c.m. angle Φ in the pd scattering and $\cos\Theta$ is as follows

$$\cos\Theta = 2 s_{pp}^{1/2} [k^2 + \frac{1}{4}(s_{pp} - m_d^2 - k^2)][(s_{pp} - m_d^2 - k^2)^2 - 4 m_d^2 k^2]^{-1/2} / |p| . \quad (4)$$

THE DWIA CROSS SECTION

In the following discussion we follow closely the paper /1/ for cluster knockout. The generalization to the $p<pn> \to pd$ elementary process is relatively simple.

A. Formulation

We use the impulse approximation within the framework of a three particle system. The $A(p,p'd)B$ reaction is described by the scheme

$$p + (B + <pn>) \to B + p + d , \quad (5)$$

where the target nucleus is considered as a complex of particles B and the <pn> subsystem.

The laboratory differential cross section for the three particle final state in the range $d\vec{k}_d$, $d\vec{k}_{p'}$, $d\vec{k}_B$, see /22/ .

$$\frac{d^9\sigma}{d\vec{k}_{p'} d\vec{k}_d \, d\vec{k}_B} = \frac{(2\pi)^4}{|\vec{v}_{rel}|} |T_{if}|^2 \delta(\sum_f E_f - \sum_i E_i)\, \delta(\sum_f \vec{k}_f - \sum_i \vec{k}_i) , \quad (6)$$

where $\vec{v}_{rel}$ is the relative velocity between the incident particle and the target, $\vec{k}_a$ and E_a are the momentum and total energy of the appropriate particle. The T_{if} amplitude, provided that the exchange terms due to the antisymmetrization between the incident proton and the residual nucleus B are neglected, has the following form:

$$T_{if} = [A(A-1)/2]^{1/2} < \Phi(\vec{B})\, \Phi(\vec{p},\vec{d}) \,|\, t_{if} \,|\, a_{p<pn>} \Phi(\vec{A})\, \Phi(p) > . \quad (7)$$

The ~ serves as a reminder that the wave functions are antisymmetrized with respect to the interchange of any two nucleons. $a_{p<pn>}$ is the antisymmetrizer between the incident proton and the $<pn>$ cluster.

The target wave function $\Phi(\tilde{A})$ can be expanded in terms of the residual nucleus B . The expansion coefficient is

$$<\Phi_{J_B M_B T_B N_B}(\tilde{B})|\Phi_{J_A M_A T_A N_A}(\tilde{A})> = \qquad (8)$$
$$= \sum_{\substack{JMt\nu \\ \alpha}} (T_B N_B t\nu | T_A N_A)(J_B M_B J M | J_A M_A)\, \Phi^{AB}_{\alpha JMt\nu}(\vec{k},\vec{P}),$$

where J_a (projection M_a), T_a (projection N_a) are the angular momentum and isospin quantum numbers, $\vec{k}$ and $\vec{P}$ represent the relative momentum between proton and neutron and the relative momentum of B with respect of the $<pn>$ center-of-mass, respectively, and other necessary quantum numbers are denoted by α . The $\Phi^{AB}_{\alpha JMt\nu}(\vec{k},\vec{P})$ function is defined by the equation (8) and can be expanded as follows:

$$\Phi^{AB}_{\alpha JMt\nu}(\vec{k},\vec{P}) = \sum_{\substack{\ell m_\ell \\ s m_s \\ J m \\ L\Lambda}} (\ell m_\ell s m_s | jm)(jmL\Lambda|JM)\, Y_{L\Lambda}(\hat{P}) Y_{\ell m_\ell}(\hat{k}) \times \qquad (9)$$
$$\times\, \chi_{sm_s} \phi^{AB}_{\alpha JLj\ell t\nu}(k,P),$$

where χ_{sm_s} is the two nucleon spin function. Using the expansion coefficient (8) and provided that t_{if} is assumed not to act upon the internal coordinates of the residual nucleus B we can integrate over these variables:

$$T_{if} = [A(A-1)/2]^{1/2} \sum_{\substack{JMt\nu \\ \alpha}} (T_B N_B t\nu | T_A N_A)(J_B M_B JM | J_A M_A) \times \qquad (10)$$
$$\times <a_{pd}\eta^{(-)}_{Bpd}\, \phi_{m_{p'}} \phi_{m_d}(\tilde{d})\, | t_{if} |\, a_{p<pn>}\eta^{(+)}_{pA}\, \phi^{AB}_{\alpha JMt\nu}(\vec{k},\vec{P})\phi_{m_p}>,$$

where $\eta^{(+)}_{pA}$ and $\eta^{(-)}_{bpd}$ describe the relative motion of mass centers of particles in the entrance and exit channels, respectively, $m_{p'}$, m_d and m_p are spin projection quantum numbers.

Introducing the impulse approximation we can replace t_{if} by the two body operator of the free $p<pn>\to pd$ reaction process t_f . We introduce the additional assumption that the amplitude of the elementary process varies sufficiently slowly with the momenta so its arguments may be replaced by their asymptotic values. This procedure leads to the zero range expression in the matrix element and T_{if} has the following form

$$T_{if}=[A(A-1)/2]^{1/2}\sum_{\substack{JMt\nu\\ \alpha}}(T_B N_B t\nu|T_A N_A)(J_B M_B JM|J_A M_A)\times$$
$$\times\int d\vec{k}<\vec{k}_p,m_p,\vec{k}_d m_d|t_f|\vec{k}_p m_p\,\vec{P}\vec{k}><\eta^{(-)}_{Bpd}|\delta(\vec{r}_p-\vec{r}_d)|\eta^{(+)}_{pA}\phi^{AB}_{\alpha JMt\nu}(\vec{k},\vec{P})>. \tag{11}$$

Now we take into account the dominance of the $\ell=0$ and $t=0$ component[20] in the $p<pn>\to pd$ process. As the radial wave functions with $\ell=0$ have their maximum at $\vec{k}=0$, by calculating the zero range matrix element at $\vec{k}=0$ we can take it out of the integral. Using formulae (9) we can write

$$T_{if}=[A(A-1)/2]^{1/2}\sum_{\substack{\alpha JM\\ L\Lambda}}(J_B M_B JM|J_A M_A)(L\Lambda 1m|JM)T^{AB}_{\alpha JL\Lambda}\times$$
$$\times<\vec{k}_p,m_p,\vec{k}_d m_d|t_f|k_p\;m_p\;\phi^{AB}_{\alpha JL}(k,P)m> \tag{12}$$

$$T^{AB}_{\alpha JL\Lambda}=\frac{(2\pi)^3}{\phi^{AB}_{\alpha JL}(k=0,P)}<\eta^{(-)}_{Bpd}\;|\delta(\vec{r}_p-\vec{r}_d)|\eta^{(+)}_{pA}\phi^{AB}_{\alpha JL}(k=0,P)Y_{L\Lambda}(\hat{P})>. \tag{13}$$

By taking the spin structure of the $pd\to pd$ amplitude from[14] in the calculation of $|T_{if}|^2$ we can sum up on the spin projection quantum number straightforwardly. For the $A(p,p'd)B$ laboratory differential cross section we can obtain the following expression.

$$\frac{d^3\sigma}{d\Omega_{p'}dT_{p'}d\Omega_d}=s_{pd}\,\frac{k_{p'}k_d}{k_p E_{NN}}\,\frac{1}{1+\dfrac{E_d}{E_B}\left(1-\dfrac{k_p}{k_d}\cos\Theta_{pd}+\dfrac{k_{p'}}{k_d}\cos\Theta_{p'd}\right)}A^2(s_{pd},u_{pd}) \tag{14}$$
$$\times A(A-1)/2\sum_{\substack{\alpha\alpha'\\ JL}}f_{\alpha JL}\;f^*_{\alpha' JL}\;D^{AB}_{\alpha\alpha' JL},$$

where

$$f_{\alpha JL} = \int_0^\infty e^{-\gamma r}(1+\gamma r)\,\phi_{\alpha JL}^{AB}(r,P)\,j_l(\tilde{p}r)\,dr \tag{15}$$

$$D_{\alpha\alpha' JL}^{AB} = \frac{1}{2L+1}\sum_\Lambda T_{\alpha JL\Lambda}^{AB}\,T_{\alpha' JL\Lambda}^{AB*} \tag{16}$$

with $E_{NN} = (m_d^2+P^2)^{1/2}$. The definition of the relativistic phase space factor is from /23/. Note that if the additional quantum numbers α are not needed the cross section formula (14) is greatly simplified.

B. Evaluation of the amplitude $T_{\alpha JL\Lambda}^{AB}$

The initial and final scattering states for the three-body system are generated using the $H_i = H_f = H - V_{pd}$ Hamiltonian. We write the Hamiltonian in the entrance channel as

$$H_i = T_{pA} + V_{pB} + \mathcal{H}_A, \tag{17}$$

where $\mathcal{H}_a$ is the internal Hamiltonian of particle a , the $V_{\alpha\beta}$ represents the interaction between particles α and β and the relative kinetic energy operator $T_{\alpha\beta}$ corresponding to the relative coordinate $\vec{r}_{\alpha\beta}$ is written as (h=1 , c=1);

$$T_{\alpha\beta} = -\frac{\vec{\nabla}_{\alpha\beta}^2}{2\mu_{\alpha\beta}}, \tag{18}$$

where $\mu_{\alpha\beta}$ is the reduced mass of α and β . The V_{pB} interaction operator is a function of the $\vec{r}_{pB}$ coordinate and not of $\vec{r}_{pA}$ and it needs carefull treatment /24/. In order to write the initial wave function in the product form we can accept the approximation that the elementary interaction V_{pd} is sufficiently short ranged so that the distortions do not change significantly over this range. One, therefore, needs the entrance and exit channel functions in the vicinity of $\vec{r}_{pd} = 0$. As

$$\vec{r}_{pA} = \frac{m_B}{m_A}\vec{r}_{pB} - \frac{m_d}{m_d+m_B}\vec{r}_{pd}, \tag{19}$$

where m_A and m_B are the target and residual nucleus masses, respectively, we can use the following approximation:

$$V_{pB}(\vec{r}_{pB}) \approx V_{pB}(\frac{m_A}{m_B}\vec{r}_{pA}) \tag{20}$$

Another possibility is the use of the $V_{pA}(\vec{r}_{pA})$ potential with both real and imaginary well depths reduced by a factor m_B/m_A [1].

In the exit channel for the Hamiltonian we have

$$H_f = T_{pB} + V_{pB} + T_{d(pB)} + V_{dB} + \mathcal{H}_d + \mathcal{H}_B , \tag{21}$$

where $T_{d(pB)}$ corresponds to the Jacobian coordinate

$$\vec{r}_{d(pB)} = \vec{r}_d - \frac{m\vec{r}_p + m_B\vec{r}_B}{m + m_B} . \tag{22}$$

The (21) form is different from the usual expression [24], where

$$H_f = T_{pB} + V_{pB} + T_{dB} + V_{dB} + \frac{\vec{\nabla}_{pB}\vec{\nabla}_{dB}}{m_B} + \mathcal{H}_d + \mathcal{H}_B . \tag{23}$$

This expression contains the $\vec{\nabla}_{pB}\vec{\nabla}_{dB}/m_B$ coupling term. Its effect in the usual coplanar-like geometry is very small but in our case where the final proton and deuteron move in opposite direction it may be very important.

In expression (22) the V_{dB} interaction depends on the $\vec{r}_{dB}$ coordinate and the appropriate kinetic energy $T_{d(pB)}$ is written in terms of the (22) $\vec{r}_{d(pB)}$ coordinate. Using like in the entrance channel the short range nature of the V_{pd} interaction we have the following expression:

$$V_{dB}(\vec{r}_{dB}) \approx V_{dB}(\frac{m+m_B}{m_B}\vec{r}_{d(pB)}) . \tag{24}$$

Solving the Schrödinger equation with the Hamiltonians (17) and (22) for the initial and final states, respectively,we obtain:

$$\eta^{(+)}_{pA} = \chi^{(+)}(\vec{k}_{pA}, \vec{r}_{pA}) , \tag{25}$$

$$\eta^{(-)}_{Bpd} = \chi^{(-)}(\vec{k}_{pB}, \vec{r}_{pB})\,\chi^{(-)}(\vec{k}_{d(pB)}, \vec{r}_{d(pB)}) , \tag{26}$$

where the $\vec{k}_{pA}$, $\vec{k}_{pB}$, $\vec{k}_{d(pB)}$ momenta are calculated from the asymptotic laboratory momenta $\vec{k}_p$, $\vec{k}_{p'}$, $\vec{k}_d$ according to the formulae:

$$\vec{k}_{pA} = \frac{m_A}{m+m_A}\vec{k}_p \,, \tag{27}$$

$$\vec{k}_{pB} = \vec{k}_{p'} - \frac{m}{m+m_B}(\vec{k}_p - \vec{k}_d), \tag{28}$$

$$\vec{k}_{d(pB)} = \vec{k}_d - \frac{m_d}{m+m_d+m_B}\vec{k}_p \,. \tag{29}$$

Substituting for $\eta^{(+)}_{pA}$ and $\eta^{(-)}_{Bpd}$ in Eq. (13) and integrating over $\vec{r}_{pd}$ we obtain

$$T^{AB}_{aJL\Lambda} = \frac{(2\pi)^3}{\phi^{AB}_{aJL}(k=0,P)} \int \chi^{(-)*}_{pB}(\vec{k}_{pB},\vec{R})\, \chi^{(-)*}_{d(pB)}(\vec{k}_{d(pB)}, a_2\vec{R}) \times$$
$$\times \chi^{(+)}_{pA}(\vec{k}_{pA}, a_1\vec{R}) \times \phi^{AB}_{aJL}(k=0,R)\, Y_{L\Lambda}(\hat{R})\, d\vec{R}, \tag{30}$$

where we have used formulae (22) and relation

$$\vec{r}_{pB} = \vec{r}_{dB} + \vec{r}_{pd} \,,$$
$$\vec{r}_{pA} = \frac{m_B}{m_d+m_B}\vec{r}_{dB} + \vec{r}_{pd} \,, \tag{31}$$
$$a_1 = m_B/(m_d+m_B), \quad a_2 = m_B/(m_p+m_B).$$

The distorted waves $\chi_a(\vec{k},\vec{r})$ are calculated according to the eikonal approximation [25]. We rewrite $\chi_a(\vec{k},\vec{r})$ as

$$\chi_a(\vec{k},\vec{r}) = D_a(\vec{r})\, e^{i\vec{k}\vec{r}} \frac{1}{(2\pi)^{3/2}} \,. \tag{32}$$

The distortion factors $D_a(\vec{r})$ are then given in a coordinate system with the origin in the center-of-mass of the residual nucleus B . Using formulae (20), (24) we get

$$D_{pA}(a_1\vec{R}) = e^{-i a_1 \frac{E_{pA}}{k_{pA}} \int_{-\infty}^{0} V_{pB}(\vec{R}+s\hat{k}_{pA})\, ds} \,, \tag{33a}$$

$$D_{pB}(\vec{R}) = e^{-i\frac{E_{pB}}{k_{pB}}\int_0^\infty V_{p'B}(\vec{R}+s\hat{k}_{pB})\,ds}, \qquad (33b)$$

$$D_{d(pB)}(\alpha_2\vec{R}) = e^{-i\alpha_2\frac{E_{d(pB)}}{k_{d(pB)}}\int_0^\infty V_{dB}(\vec{R}+s\hat{k}_{d(pB)})\,ds}, \qquad (33c)$$

where E_α are the total energies corresponding to the momenta k_α. The $V_{pB}(r)$, $V_{p'B}(r)$ potentials are the proton-nucleus central, complex optical potentials calculated from the isospin averaged forward nucleon-nucleon scattering amplitudes at the kinetic energies corresponding to the momenta $\vec{k}_{pB}$, $\vec{k}_{p'B}$, respectively, following the method of ref.[26]. The form of the deuteron-nucleus optical potential $V_{dB}(r)$ can be constructed from the nucleon-nucleus optical potentials by the usual folding procedure[27]. For the final formulae of the amplitude $T^{AB}_{\alpha JL\Lambda}$ we get

$$T^{AB}_{\alpha JL\Lambda} = \frac{1}{(2\pi)^{3/2}\phi^{AB}_{\alpha JL}(k=0,P)}\int e^{i\vec{P}\vec{R}}\, D^*_{pB}(R)\, D^*_{d(pB)}(\alpha_2\vec{R})\, D_{pA}(\alpha_1\vec{R})\times \qquad (34)$$

$$\times\, \phi^{AB}_{\alpha JL}(k=0,R)\, Y_{L\Lambda}(\hat{R})\, d\vec{R}$$

where we have made use of the (27)-(29) and (32) relations.

THE ANALYSIS OF THE $^6Li(p,p'd)^4He$ QUASIFREE SCATTERING

The $^6Li(p,p'd)^4He$ quasifree large-angle scattering was investigated in a kinematically complete experiment at 670 MeV incident energie[28]. The forward deuteron angle was chosen to be 6.5° the backward proton angle region was taken from -140 to -152° in the laboratory system. The overall energy resolution was ~17 MeV which made possible the separation of events leading to the ground state of residual nucleus from the events leading to excited states. The analysis of experiment gives the impossible effective number of deuterons in the p-shell $N_{eff} = 1.08 \pm 0.1$ which means the failure of the model constructed in terms of the overlap integral.

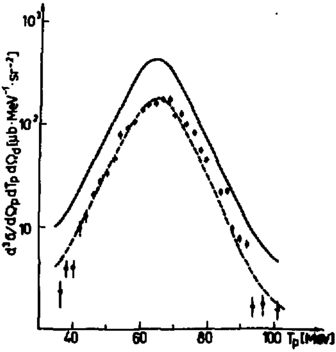

Fig.2. The ^{6}Li(p,p′d)^{4}He differential cross section at T_p = 670 MeV, Θ_d = 6.5°, $\Theta_{p'}$ = −147° due to the three-body wave function Rai, Lehman and Ghovanlou (solid line) and including short range correlations (dotted line). Experimental points are taken from /28/.

The relatively simple structure of the nucleus ^{6}Li gives a good opportunity for the application of our model. The ground state of ^{6}Li can be regarded as an α−p−n three-body system, if we neglect the structure of the α-particle. For the description of the α−p−n system we have used a three-body wave function of ref./10/. They have solved the Faddeev equations using for the n−p and N−p interactions separable potentials. The n−p interaction is taken to be $^3S_1$, and the potential to be of a standard Yamaguchi form.

The three-body wave function normalized to unity has the form $\psi(\vec{k},\vec{P})$, where $\vec{k}$ and $\vec{P}$ represent the relative momentum between proton and neutron and the relative momentum of ^{4}He with respect to the p−n center of mass, respectively. Now in the cross section the index α is not needed and the main contribution is given by the J = 1, L = 0 component.

The proton − ^{4}He optical potential for the distortion factors are calculated assuming Gaussian form for the density function of ^{4}He.

The ^{6}Li(p,p′d)^{4}He differential cross section at incident proton energy 670 MeV at forward deuteron angle 6.5° and backward proton angle Θ_p = −147° as a function of the backward proton kinetic energy T_p is presented in fig.2. It can be seen that the calculated cross section (solid line) is about twice as much as the experimental values of ref./28/. We notice that the calculation without distortion, in plane wave approximation gives a cross section with a 30% higher values. Taking into account the relatively small effect of the distortion the approximations made in its calculation do not influence essentially the calculated cross section. The contribution of the L = 2 component is less than 1 per cent.

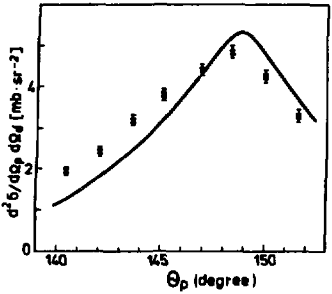

Fig.3. The $^6Li(p,p'd)^4He$ proton angular distribution at T_p = 670 MeV, Θ_d = 6.5° due to the three-body wave function including correlations. Experimental points are taken from /28/.

The difference between the experimental and theoretical values can be explained in the following manner. The simple Yamaguchi factor which represents the proton-neutron interaction does not contain the repulsive core appearing in the nucleon-nucleon interaction. The short range repulsion can be taken into account by including into the $\psi(\vec{r},\vec{P})$ wave function the following correlation function/29/:

$$f(r) \approx 0\,, \quad r < r_c\,,$$
$$f(r) = 1 - e^{-\beta(r/r_c - 1)}, r \geq r_c\,. \qquad (35)$$

If we chose the parametrization r_c = 0.4 fm, β = 1.5 then the calculated cross section (dotted line) decreases by half and we have a satisfactory agreement with the experiment. The above values of parameters r_c and β correspond to a hard core in the deuteron wave function.

The angular distribution of the backward protons averaged over the proton energies we presented in Fig.3. For the calculation we have used the three-body wave function multiplied by the correlation function with the above parameters r_c and β. We have a satisfactory agreement with the experimental angular distribution.

SUMMARY

In the theory of the $A(p,p'd)B$ quasifree scattering we have to study carefully the character of the $p\langle pn\rangle \to pd$ elementary process as the exchange and other effects depending directly on the inner structure of the $\langle pn\rangle$ cluster may have an important role. In the latter case the concept of the effective number of deuterons in the target nucleus is useless and misleading.

In the case of the $A(p,p'd)B$ quasifree large angle scattering at intermediate incident energies the $p\langle pn\rangle \to pd$ elementary process can be described by the triangle mechanism of Craigie and Wilkin. The $p\langle pn\rangle \to pd$ amplitude depends only on the short range part of the $\langle pn\rangle$ wave function. The DWIA cross section can be formulated in a relatively simple form, the distortion is calculated in eikonal approximation. The possibilities of the model are demonstrated in the analysis of the $^6Li(p,p'd)^4He$ scattering.

Taking into account the strong absorption of the deuteron and proton in the final state the method can be applied for the study of $p-n$ pairs on the surface of the nucleus. The model can be easily extended for the investigation of the $\langle nn\rangle$ and $\langle pp\rangle$ pairs /20/. As the main contribution to the cross section is due to the nucleon pairs with $\ell = 0$ relative orbital momentum where the two nucleons are very close to each other, the experiments can give useful information on the pairing effects on the surface of the nuclei.

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