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メタデータ

| 言語   | eng |
|--------|-----|
| 出版者  |     |
| 公開日  | 2008-01-16 |
| キーワード (Ja): |     |
| キーワード (En): |     |
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URL http://hdl.handle.net/10228/691
Spin dependent momentum distributions of proton-deuteron clusters in $^3$He from electron scattering on polarized $^3$He: Theoretical predictions

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(Received 21 February 2002; published 28 May 2002)

The process $^3$He($e,e'p)d$ or $^3$He($e,e'd)p$ is studied theoretically in a Faddeev treatment with the aim to have access to the spin dependent momentum distribution of $p\bar{d}$ clusters in polarized $^3$He. Final state interactions and meson exchange currents turn out to have a strong influence in the considered kinematical regime (below the pion threshold). This precludes direct access to the momentum distribution except for small deuteron momenta. Nevertheless, the results for the longitudinal and transverse response functions are interesting as they reflect our present day understanding of the reaction mechanism and therefore data would be very useful.

DOI: 10.1103/PhysRevC.65.064004 PACS number(s): 24.70.+s, 25.10.+s, 25.40.Lw

I. INTRODUCTION

With knowledge of solving precisely few-nucleon equations, the availability of high-precision nucleon-nucleon ($NN$) potentials and insight into the electromagnetic nucleonic current operator it is seducing to ask very detailed questions about spin dependent momentum distributions inside light nuclei and the way to access them through electron scattering taking final state interactions fully into account. Momentum distributions of polarized $\vec{p}\vec{\bar{d}}$ clusters in spin-oriented $^3$He have been studied before; see, for instance, [1]. We address here the question whether these distributions are accessible through the $^3$He($e,e'p)d$ or $^3$He($e,e'd)p$ processes. Optimal kinematical conditions are that the polarizations of $^3$He and of the knocked out proton (deuteron) and the momenta of the final proton and deuteron are collinear to the photon momentum. As we will show the longitudinal and transverse response functions will lead, up to known factors, directly to the sought spin dependent momentum distribution of the $\vec{p}\vec{\bar{d}}$ clusters in $^3$He. One can also define a proper asymmetry, which carries corresponding information. Of course this can only be true in a plane-wave impulse approximation (PWIA) and for the absorption of the photon on a single nucleon. Rescattering effects in the final state as well as meson exchange currents (MECs) will disturb the outcome. The strength of that disturbance again will depend on the photon momentum $Q$ with the hope that it decreases with increasing $Q$.

We formulate the electromagnetic process in Sec. II and also display there the $\vec{p}\vec{\bar{d}}$ cluster momentum distributions of $^3$He. Section III shows our results for the $^3$He($e,e'p)d$ and $^3$He($e,e'd)p$ processes based on the AV18 $NN$ potential [2] and precise solutions of the corresponding Faddeev equations. Since our predictions depend on the full dynamics in a highly nontrivial manner, a future experimental verification will be an important test for the understanding of few-nucleon dynamics. We end with a brief summary in Sec. IV.

II. THEORY

The spin dependent momentum distribution of proton-deuteron clusters inside the $^3$He nucleus is defined as

$$\mathcal{V}(M,M_d,m;\vec{q}_0) = \langle \Psi M \mid \phi_d M_d \mid \vec{q}_0 \frac{1}{2} m \rangle \langle \vec{q}_0 \frac{1}{2} m \mid \phi_d M_d \mid \Psi M \rangle,$$

(1)

where $\vec{q}_0$ is the proton momentum (the deuteron momentum is $-\vec{q}_0$); $m, M_d$, and $M$ are spin magnetic quantum numbers for the proton, deuteron, and the considered nucleus, respectively.

We introduce our standard basis in momentum space [3]

$$|pq\alpha\rangle = pq(\lambda sj) \left( \begin{array}{c} 1 \end{array} \right) |jM\rangle \left( \begin{array}{c} 1 \end{array} \right) T_M,$$

(2)

where $p$ and $q$ are magnitudes of Jacobi momenta and the set of discrete quantum numbers $\alpha$ comprises angular momenta, spins, and isospins for a three-nucleon (3N) system. Then $\mathcal{V}(M,M_d,m;\vec{q}_0)$ can be evaluated as

$$\mathcal{V}(M,M_d,m;\vec{q}_0) = \sum_a \left( \delta_{l_0} + \delta_{l_2} \right) \delta_{s_1} \delta_{j_1} \delta_{t_0} \left| \frac{1}{2} j M_d, M - M_d, M \right| \right)^2 \times \int_0^\infty dp \, p^2 \phi(p) \langle pq_0 | \Psi \rangle Y_{l_0,j_0}^* (\vec{q}_0) \right| \frac{1}{2} j M_d, M - M_d, M \right|$$

(3)
FIG. 1. Absolute value of $H_\lambda(q_0)$ defined in Eq. (5) for $\lambda=0$ (solid line) and $\lambda=2$ (dashed line). Note $H_0(q_0)<0$ for $q_0\geq 400$ MeV/c, while $H_2(q_0)$ remains always positive for the shown $q_0$ values.

In Eq. (3), $\langle p q_0 \alpha | \Psi \rangle$ are the partial-wave projected wave function components of $^3$He in momentum space and $\phi_\ell(p)$ are the $s$- and $d$-wave components of the deuteron.

Further we rewrite $\mathcal{Y}(M, M_d, m; \vec{q}_0)$ as

$$\mathcal{Y}(M, M_d, m; \vec{q}_0) = \left| \sum_{\lambda=0,2} Y_{\lambda, M-M_d-m}(\vec{q}_0) C \left( I_{\lambda} \frac{1}{2}, M-M_d, M \right) \times C \left( \frac{1}{2} I_{\lambda} \frac{1}{2}, M-M_d, m, M-M_d \right) \right|^2$$

and define an auxiliary quantity $H_\lambda(q_0)$ as

$$H_\lambda(q_0) = \sum_{l=0,2} \int_0^\infty dp \ p^2 \phi_\ell(p) \langle p q_0 \alpha | \Psi \rangle.$$

Note that the set $\alpha_\lambda$ contributes only for the deuteron quantum numbers $s=1, j=1,$ and $t=0$. Further $I_{\frac{1}{2}}$ for $\lambda$

and shown in Fig. 3.

Next we ask the question how this quantity can be accessed experimentally. The cross section for the process $e^+ + ^3$He→$e^+ + p + d$ has the form [4]

$$\sigma = \sigma_{\text{Mott}} (u_L W_L + v_T W_T + \nu_{TT} W_{TT} + \nu_{TL} W_{TL})$$

$$+ h (v_T W_T + \nu_{TL} W_{TL}) \rho,$$

where $\sigma_{\text{Mott}}, v_i,$ and $\rho$ are analytically given kinematical factors, and $h$ is the helicity of the incoming electron. The response functions $W_i$, which contain the whole dynamical

information, are constructed from the current matrix elements taken between the initial bound state $| \Psi_M \rangle$ and the final scattering state $| \Psi_{pd}^{s} M_d m \rangle$ [5]. They are given as

$$W_L = | \langle \Psi_{pd}^{s} M_d m | j_0(\vec{Q}) | \Psi_M \rangle |^2 = |N_0|^2,$$

$$W_T = | \langle \Psi_{pd}^{s} M_d m | j_1(\vec{Q}) | \Psi_M \rangle |^2$$

$$+ | \langle \Psi_{pd}^{s} M_d m | j_{-1}(\vec{Q}) | \Psi_M \rangle |^2 = |N_+|^2 + |N_-|^2,$$

$$W_{TT} = 2 \text{Re}[N_+(N_-)^*].$$
FIG. 3. The asymmetry $A = [\langle m = \frac{1}{2} \rangle - \langle m = - \frac{1}{2} \rangle] / [\langle m = \frac{1}{2} \rangle + \langle m = - \frac{1}{2} \rangle]$. 

$W_{TL} = -2 \text{Re}[N_0(N_+ - N_-)\ast]$. 

$W_T = |N_+|^2 - |N_-|^2$. 

$W_{TL'} = -2 \text{Re}[N_0(N_+ + N_-)\ast]$. 

Note that $W_T$ and $W_{TL'}$ contribute only in the case when the initial electron is polarized. This is our standard notation $N$ of the nuclear matrix element, where the indices 0 and ±1 stand for the zeroth component and the transverse spherical components of the current. The general $3N$ current operator contains the single-nucleon contributions as well as two- and three-nucleon exchange terms

\[ j_\mu(\hat{Q}) = j_\mu(\hat{Q}; 1) + j_\mu(\hat{Q}; 2) + j_\mu(\hat{Q}; 3). \] 

In the nonrelativistic limit, which we use, the three contributing pieces of the single-nucleon current operator (the charge density, the convection, and the spin current) can be written in the $3N$ momentum space as

\[ j_\mu(\hat{Q}; 1) = \int d\vec{p} d\vec{q} \langle \hat{p} \hat{q} | -\frac{2}{3} \hat{Q} \rangle \Pi(Q) \left| -\frac{2}{3} \hat{Q} \right|, \]

\[ j_\mu(\hat{Q}; 2; \text{conv}) = \int d\vec{p} d\vec{q} \langle \hat{p} \hat{q} | \frac{q_\sigma}{m_N} \Pi(Q) \left| -\frac{2}{3} \hat{Q} \right|, \]

\[ j_\mu(\hat{Q}; 3; \text{spin}) = \int d\vec{p} d\vec{q} \langle \hat{p} \hat{q} | \frac{Q \sigma \tau}{2m_N} \Pi(Q) \left| -\frac{2}{3} \hat{Q} \right|. \]

where $m_N$ is the nucleon mass and $\Pi(Q)$ and $\Pi_M(Q)$ are sums of isospin projection operators for the neutron and proton joined by the electric ($G_E$) and magnetic ($G_M$) nucleon form factors, respectively (see [5]). We assumed that $\hat{Q} \parallel \hat{z}$.

Let us now decompose the scattering state $|\Psi_{pd}^{(-)}M_{d,m}\rangle$ in the following way:

\[ |\Psi_{pd}^{(-)}M_{d,m}\rangle = |\phi_{d}M_{d,\uparrow}\rangle + |\Psi_{pd}^{\ast}M_{d,m}\rangle. \]

The first term is just a product of the deuteron wave function $|\phi_{d}M_{d,\uparrow}\rangle$ and a relative momentum eigenstate of the spectator nucleon $|\tilde{q}_{12}m\rangle$. The other term accounts for the proper antisymmetrization of the final state and all rescattering contributions.

If the many-nucleon contributions to the $3N$ current $[j_\mu(\hat{Q}; 2) \text{ and } j_\mu(\hat{Q}; 3)]$ and $|\Psi_{pd}^{\ast}M_{d,m}\rangle$ can be neglected (PWIA assumption), then the current matrix elements take the following form:

\[ \begin{align*}
N_{0}^{\text{PWIA}}(M, M_d, m) &= G_E(Q) \sum_{\alpha} \left( \delta_{10} + \delta_{12} \right) \delta_{\alpha \beta} \delta_{\beta \alpha} C \left( \frac{1}{2}; M_d, M - M_d, M \right) C \left( \frac{1}{2}; M - M_d - m, m, M - M_d \right) \\
&\times Y_{\lambda, M - M_d - m} \left( \frac{-2}{3} \hat{Q} \right) \int_{0}^{\infty} dp \, p^2 \langle \tilde{q} \hat{q} - \frac{2}{3} \hat{Q} | \alpha | \Psi \rangle \phi_{\lambda}(p), \\
N_{\tau}^{\text{conv PWIA}}(M, M_d, m) &= \frac{\sqrt{3} \pi \hat{q} \hat{q}}{m_N} Y_{1, \lambda}(\hat{q} \hat{q}) N_{0}^{\text{PWIA}}(M, M_d, m), \\
N_{\tau}^{\text{spin PWIA}}(M, M_d, m) &= \frac{\sqrt{3}}{2} \frac{Q}{m_N} G_M(Q) \sum_{\alpha} \left( \delta_{10} + \delta_{12} \right) \delta_{\alpha \beta} \delta_{\beta \alpha} C \left( \frac{1}{2}; M_d, M - M_d, M \right) \\
&\times C \left( \frac{1}{2}; M - M_d - m + \tau, m, \tau, m - M_d \right) Y_{\lambda, M - M_d - m + \tau} \left( \frac{-2}{3} \hat{Q} \right) \\
&\times \int_{0}^{\infty} dp \, p^2 \langle \tilde{q} \hat{q} - \frac{2}{3} \hat{Q} | \alpha | \Psi \rangle \phi_{\lambda}(p).
\end{align*} \]
In the laboratory frame \( \vec{p}_N + \vec{p}_d = \vec{Q} \) and by definition of the Jacobi momentum \( \vec{q}_f = \frac{1}{2} \vec{p}_N - \frac{1}{3} \vec{p}_d \), thus \( \vec{q}_f - \frac{1}{3} \vec{Q} = - \vec{p}_d \). The second argument of the \( ^3 \)He wave function component is therefore just the deuteron laboratory momentum. For the parallel kinematics \( (\vec{Q} || \vec{p}_N || \vec{p}_d) \) the matrix element \( N_{\text{conv PWIA}} \) is zero.

In this particular situation and for the initial target spin parallel to \( \vec{Q} \) \( (M = \frac{1}{2}) \) only few combinations of the magnetic quantum numbers contribute to the nuclear matrix elements \( N^0_{\text{PWIA}} \) and \( N^\pm_{\text{spin PWIA}} \). Because of the choice of the parallel kinematics and the property of the spherical harmonics these are \( M = \frac{1}{2}, M_d = 0, m = \frac{1}{2} \) and \( M = \frac{1}{2}, M_d = 1, m = - \frac{1}{2} \) in \( N^0_{\text{PWIA}} \), \( M = \frac{1}{2}, M_d = 0, m = - \frac{1}{2} \) and \( M = \frac{1}{2}, M_d = - 1, m = \frac{1}{2} \) in \( N^1_{\text{spin PWIA}} \), and \( M = \frac{1}{2}, M_d = 1, m = \frac{1}{2} \) in \( N^1_{\text{spin PWIA}} \).

Furthermore, if we compare the expressions given in Eqs. (14) and (16) to the one in Eq. (3), we find that the spin dependent momentum distributions \( \mathcal{Y} \) of \( ^3 \)He are connected to \( N^1_{\text{PWIA}} \) by

\[
\mathcal{Y} \left( M = \frac{1}{2}, M_d = 0, \ m = \frac{1}{2}; |\vec{p}_d| \right) = \frac{1}{(G_E)^2} \left| N^0_{\text{PWIA}} \left( M = \frac{1}{2}, M_d = 0, \ m = \frac{1}{2} \right) \right|^2
\]

\[
= \frac{2m_N^2}{Q^2 (G_M)^2} \left| N^1_{\text{spin PWIA}} \left( M = \frac{1}{2}, M_d = 0, \ m = - \frac{1}{2} \right) \right|^2
\]

(17)

and by

\[
\mathcal{Y} \left( M = \frac{1}{2}, M_d = 1, \ m = - \frac{1}{2}; |\vec{p}_d| \right) = \frac{1}{(G_E)^2} \left| N^0_{\text{PWIA}} \left( M = \frac{1}{2}, M_d = 1, \ m = - \frac{1}{2} \right) \right|^2
\]

\[
= \frac{2m_N^2}{Q^2 (G_M)^2} \left| N^0_{\text{spin PWIA}} \left( M = \frac{1}{2}, M_d = 1, \ m = \frac{1}{2} \right) \right|^2
\]

(18)

In the case of parallel kinematics \( W_{TT}, W_{TL}, \) and \( W_{TL} \), vanish. This follows from the fact that the conditions on the magnetic quantum numbers, \( M, M_d, \) and \( m \), given in products of \( N_0, N^{+1}, \) and \( N^{-1} \), cannot be simultaneously fulfilled. For an experiment with unpolarized electrons, the cross section (7) contains then only the longitudinal \( (W_L) \) and transverse \( (W_T) \) response functions:

\[
\sigma = \sigma_{\text{Mot}}(v_L W_L + v_T W_T) \rho.
\]

Thus the standard “L-T” separation is required in order to access individually \( W_L \) and \( W_T \).

Another possibility is offered by an experiment with a polarized electron beam. In this case no further separation of response functions is required, since

\[
\frac{1}{2} \left[ \sigma(h = +1) - \sigma(h = -1) \right] \frac{1}{v_T \rho} = |N_+|^2 - |N_-|^2.
\]

(20)

Therefore under these extreme simplifying assumptions the response functions \( W_L, W_T, \) and \( W_{LT} \), carry directly the desired information. Note that in case of \( W_T (W_{LT}) \) only one of the two parts gives a nonzero contribution.

The full dynamics adds antisymmetrization in the final state. [Note our single nucleon current operator as given in Eqs. (10)–(12) acts only on one particle. Antisymmetrization in the final state is equivalent to the action of the current on all three particles.] Then of course rescattering to all orders in the \( NN t \) operator has to be included. On top one should add at least two-body currents. We have described how to do that before at several places [5]. Here we only remark that we employ standard \( \pi \) - and \( \rho \) -like exchange currents related to the \( NN \) force AV18, which we use throughout the paper, and that adequate Faddeev equations for \( ^3 \)He and for the treatment of FSI have been solved precisely.

### III. RESULTS

Since we work strictly nonrelativistically we want to keep the \( 3N \) c.m. energy \( E_{3N}^{c.m.} \) below the pion threshold. But in that regime we would like to study many kinematical configurations and also include higher three-momenta of the photon. We display in Table I the kinematical conditions, for which our studies have been carried through. In parallel kinematics one can distinguish three cases for the momentum orientations of the final proton and deuteron, which we denote by \( C_1, C_2, \) and \( C_3 \), and which are depicted in Fig. 4. Thus for \( C_2 \) the final momenta of proton and deuteron are parallel to \( \vec{Q} \), whereas in \( C_1 \) and \( C_3 \) only one of them lies in the direction of \( \vec{Q} \), the other is opposite. Table I shows for an (arbitrarily selected) initial electron energy of 1.2 GeV various relevant variables: the electron scattering angle, the proton and deuteron momenta \( p_N \) and \( p_d \), the photon energy \( \omega \), the three-momentum of the photon \( \vec{Q} \), and finally the \( 3N \) c.m. energy \( E_{3N}^{c.m.} \). The additional label distinguishes the three cases \( C_1, C_2, \) and \( C_3 \). We see that for each fixed \( p_d \) value we cover a certain range of \( Q \) values. The three \( C_1 \) configurations with \( E_{3N}^{c.m.} \gtrsim 140 \text{ MeV} \) are above the pion threshold and have to be taken with caution. We evaluated all the cases of Table I but do not show all in case the results are similar. Figure 5 displays \( W_L(G_E)^2 \) for \( M_d = 0, \ m = \frac{1}{2} \) and \( M_d = - \frac{1}{2} \) against the available \( Q \) values according to Table I. According to Eqs. (17) and (18), in the PWIA, \( W_L(G_E)^2 \) is just the sought \( \mathcal{Y} \) and thus trivially independent of \( Q \). Symmetrizing the final state but still neglecting rescattering is called PWIAS, while predictions including addition-
TABLE I. Electron kinematics together with different kinematical quantities used to extract the spin dependent momentum distributions of proton-deuteron clusters in $^3$He.

| $\theta_c$ (deg) | $p_N$ (MeV/c) | $p_d$ (MeV/c) | $\omega$ (MeV) | $Q$ (MeV/c) | $E^{cm}_{SN}$ (MeV) |
|----------------|---------------|---------------|----------------|-------------|---------------------|
| 14.45          | 310           | 10            | 56.67          | 300         | 35.22               |
| 19.43          | 410           | 10            | 95.01          | 400         | 61.14               |
| 24.56          | 510           | 10            | 144.00         | 500         | 94.15               |
| 29.91          | 610           | 10            | 203.63         | 600         | 134.26              |
| 35.58          | 710           | 10            | 273.92         | 700         | 181.47              |
| 14.21          | 400           | 100           | 93.33          | 300         | 71.89               |
| 19.11          | 500           | 100           | 141.26         | 400         | 107.38              |
| 24.15          | 600           | 100           | 199.83         | 500         | 149.98              |
| 29.39          | 700           | 100           | 269.05         | 600         | 199.68              |
| 19.41          | 200           | 200           | 37.44          | 400         | 3.56                |
| 24.52          | 300           | 200           | 64.06          | 500         | 14.21               |
| 29.85          | 400           | 200           | 101.33         | 600         | 31.96               |
| 35.46          | 500           | 200           | 149.25         | 700         | 56.81               |
| 41.46          | 600           | 200           | 207.82         | 800         | 88.76               |
| 16.93          | 50            | 300           | 30.80          | 350         | 3.58                |
| 27.05          | 250           | 300           | 62.74          | 550         | 3.58                |
| 29.70          | 300           | 300           | 77.39          | 600         | 8.02                |
| 35.22          | 400           | 300           | 114.66         | 700         | 22.21               |
| 41.10          | 500           | 300           | 162.58         | 800         | 43.51               |
| 21.94          | 50            | 400           | 49.45          | 450         | 8.04                |
| 35.06          | 300           | 400           | 96.04          | 700         | 3.60                |
| 40.80          | 400           | 400           | 133.32         | 800         | 14.25               |
| 14.21          | 200           | 500           | 93.41          | 300         | 71.96               |
| 19.47          | 100           | 500           | 77.44          | 400         | 43.56               |
| 24.05          | 10            | 500           | 72.17          | 490         | 24.08               |
| 13.31          | 300           | 600           | 149.35         | 300         | 127.90              |
| 19.28          | 200           | 600           | 122.73         | 400         | 88.86               |
| 24.62          | 100           | 600           | 106.76         | 500         | 56.91               |
| 29.32          | 10            | 600           | 101.49         | 590         | 34.23               |

and then to $C_3$. Symmetrization (PWIAS) has little effect at $p_d=10$ (not shown) and 100 MeV/c but has a big one for the smaller $Q$ values in case of $p_d=200–400$ MeV/c and for all $Q$ values in case of $p_d=500–600$ MeV/c. Recattering plays mostly a strong role. In the case of $C_1$ ($p_d=10$ and 100 MeV/c) its effects are relatively small and diminish nicely with increasing $Q$. In the case of $C_2$ ($p_d=200–400$ MeV/c) its role is dramatic for $p_d=300$ and 400 MeV/c, which has to be expected since the proton and the deuteron travel together with a low relative energy $E^{cm}_{SN}$. In the case of $C_3$ the two particles travel again opposite to each other as for $C_1$ and $E^{cm}_{SN}$ decreases with increasing $Q$. In this case the by-far dominant contribution to the very strong deviation from the PWIA comes from antisymmetrization in the final state and FSIs leads to a relatively mild modification in case of $m=\frac{1}{2}$ but a significantly larger one for $m=-\frac{1}{2}$. Thus we see quite different outcomes depending on the cases and these theoretical predictions would be very interesting to be compared to data.

In case of $W_T$ and $W_{TT}$ the spin operator appears in the current and moreover one can see the effects of the $\tau$- and $\rho$-like MECs. Nevertheless, the situation for $[2m_d^2/Q^2(GM)^2]W_T$ shown in Fig. 6 is roughly spoken similar to the one for $W_L/(G_E)^2$. (We regard only $W_T$ but of course $W_{TT}$ carries the same information.) Additionally one observes the effects of MECs, which are pronounced for $p_d=300$ and 400 MeV/c.

In view of all that, can we identify kinematic regions to pin down the spin dependent momentum distributions using $W_L$ or $W_T$? We choose the cases of the closest approach of $p_d$ to 300 MeV/c.

FIG. 4. Three-momenta arrangements $C_1$, $C_2$, and $C_3$ for parallel kinematics. See Table I.
FIG. 6. $[2m^2/Q^2(G_M^2)]W_T$ as a function of the three-momentum transfer $Q$ for different $p_d$ values. The curves correspond to PWIA (dotted line), PWIAS (dashed line), Full without MEC (dash-dotted line), and Full including MEC (solid line) results. The thick curves are for the $M_\perp=\frac{1}{2}$, $M_\parallel=0$, $m=\frac{1}{2}$ case, the thin lines for the $M_\perp=\frac{1}{2}$, $M_\parallel=1$, $m=\frac{1}{2}$ combination of the spin magnetic quantum numbers. In case of $p_d=400$ MeV/c the two PWIA results overlap.

PWIA and “Full” calculations (with MECs in case of $W_T$) for the different $p_d$ values. They are displayed in Figs. 7 and 8 together with the spin dependent momentum distributions $\mathcal{Y}$ from Fig. 2. In case of $m=1/2$ the values of closest approach extracted from $W_L$ and $W_T$ differ for the larger $q_0$ values, where they also do not reach $\mathcal{Y}$. Only to the left of the zero of $\mathcal{Y}$ do they agree with each other and with $\mathcal{Y}$. For $m=-1/2$ the predictions for $W_L$ and $W_T$ agree with each other but do not show the strong dip of $\mathcal{Y}$. For the smaller $q_0$ values they agree with $\mathcal{Y}$. As a consequence of these results the asymmetry $A$ formed out of those values of closest approach cannot follow the asymmetry formed out of the $\mathcal{Y}$’s. Only the values extracted for $W_L$ show a mild similarity with the asymmetry $A$, as shown in Fig. 9.

FIG. 7. $\mathcal{Y}(M=\frac{1}{2}, M_\parallel=0, m=\frac{1}{2}; q_0)$ (solid curve) as a function of the relative proton-deuteron momentum $q_0$ together with the values of closest approach (see text) from $W_L$ (squares) and from $W_T$ (circles).

FIG. 8. $\mathcal{Y}(M=\frac{1}{2}, M_\parallel=1, m=\frac{1}{2}; q_0)$ (solid curve) as a function of the relative proton-deuteron momentum $q_0$ together with the values of closest approach (see text) from $W_L$ (squares) and from $W_T$ (circles).

FIG. 9. The asymmetry $A=[\mathcal{Y}(m=\frac{1}{2})-\mathcal{Y}(m=\frac{1}{2})]/[\mathcal{Y}(m=\frac{1}{2})+\mathcal{Y}(m=-\frac{1}{2})]$ as a function of the relative proton-deuteron momentum $q_0$ together with the values of closest approach (see text) from $W_L$ (squares) and from $W_T$ (circles).

IV. SUMMARY

Based on the $NN$ force AV18 and consistent $\pi$- and $\rho$-like exchange currents we investigated within the Faddeev framework the process $^3\text{He}(e,e'p)d$ [or $^3\text{He}(e,e'd)p$]. The aim was to have access to the spin dependent momentum distribution of polarized $pd$ clusters in polarized $^3\text{He}$. That distribution would provide interesting insight into the $^3\text{He}$ wave function. We restricted ourselves to a nonrelativistic regime, where the $3N$ c.m. energy of the final state should stay below the pion threshold. In that kinematical regime we explored the longitudinal and transverse response functions $W_L$ and $W_T$, as well as $W_T$, as a function of the final deuteron and the allowed photon momenta. All the spins and momenta are chosen parallel or antiparallel to the photon momentum. While in the PWIA $W_L$ and $W_T$ ($W_T$) up to known factors yield directly the sought spin dependent momentum distribution, FSIs and MECs preclude in most cases the direct access to that distribution. The response functions $W_L$ and $W_T$ multiplied by appropriate factors have been mapped out in a wide kinematical range and this theoretical outcome should be checked experimentally. It presents the present day state-of-the-art insight into the dominant photon absorption process and the few-nucleon dynamics. It is only at small deuteron momenta $p_d=2$ fm$^{-1}$ that the investigated momentum distribution can be accessed within the constrained kinematics we have chosen.

Right now we have no reliable estimate for the amount of relativistic corrections or insight into the stability of our results under exchange of nuclear forces and consistent MECs. Clearly work in that respect should be envisaged.
ACKNOWLEDGMENTS

We are indebted to Dr. H. Gao and Dr. D. Dutta, who inspired us to perform this study. We would also like to thank Dr. M. O. Distler for an interesting discussion about experiments with polarized electrons. This work was supported by the Deutsche Forschungsgemeinschaft (J.G.), the Polish Committee for Scientific Research under Grant Nos. 2P03B02818 and 2P03B05622, and by NFS Grant No. PHY0070858. R.S. acknowledges the Foundation of Polish Science for financial support. The numerical calculations have been performed on the Cray T90 of the NIC in Jülich, Germany.

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