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Study of deuteron-proton charge exchange reaction at small transfer momentum

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Abstract. The charge exchange reaction $pd \to npp$ at 1 GeV projectile proton energy is studied in the multiple-scattering expansion technique. This reaction is considered in a special kinematics, when the momentum transfer from the beam proton to the fast neutron is close to zero. The differential cross-section and a set of polarization observables are calculated. It was shown that the contribution of the final-state interaction between two protons is very significant.

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1 Introduction

During the last decades the deuteron-proton charge exchange reaction has been studied both from the experimental and the theoretical point of view. The considerable interest in this reaction is connected, first of all, to the opportunity to extract some information about the spin-dependent part of the elementary nucleon-nucleon charge exchange amplitudes. This idea was suggested by Pomeranchuk [1] already in 1951, but until now it continues to be of interest. Later, this supposition has been developed in [2–4]. It was shown, that in the plane-wave impulse approximation (PWIA) the differential cross-section and tensor analyzing power T_{20} in the dp charge exchange reaction are actually fully determined by the spin-dependent part of the elementary $np \to pn$ amplitudes. Analogous result was obtained in [5], where this process has been studied in the Bethe-Salpeter formalism.

The differential cross-section of the $dp \to npp$ reaction at 3.34 GeV/c deuteron beam has been measured in 1970's in the 1m hydrogen bubble chamber of the JINR Synchrophasatron [6,7]. However, the obtained statistics is not sufficient to evaluate the magnitude of the spin-dependent part of the elementary amplitude. Nowadays the experiment on the study of the dp charge exchange reaction at the small momentum transfer in the GeV-region is planned at the ANKE setup at COSY [8]. The aim of this experiment is to provide information about spin-dependent np elastic-scattering amplitudes in the energy region where phase-shift analysis data are absent.

From our point of view, under the kinematical conditions proposed in this experiment, when the momentum of the emitted neutron has the same direction and magnitude as the beam proton (in the deuteron rest frame), and the relative momentum of the two protons is very small, the final-state interaction (FSI) effects play a very important role. The contribution of the D-wave in the DWF to the differential cross-section in this kinematics must be negligible [9]. However, for the polarization observables the influence of the D-component can be significant.

The goal of our paper is to study the importance of the D-wave and FSI effects under the kinematical conditions of the planned experiment. We consider the $pd \to npp$ reaction in the approach which has been used by us to describe the pd breakup process at 1 GeV projectile proton energy [10]. This approach is based on the Alt-Grassberger-Sandhas formulation of the multiple-scattering theory for the three-nucleon system. The matrix inversion method [11,12] has been applied to take account of the FSI contributions. Since the unpolarized and polarized mode of the deuteron beam are supposed to be employed in the experiment, we also calculate both the differential cross-section and a set of polarization observables. It should be noted that in this paper we have not considered the Coulomb interaction in the pp-pair. This problem is nontrivial and requires a special investigation.

The paper is organized as follows. In sect. 2 a short description of the general theoretical formalism is given. The special kinematics, when the momentum transfer from the

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beam proton to the neutron is close to zero, is considered in sect. 3. The results of our calculations for the differential cross-section and polarization observables are presented in sect. 4. The figures in this section demonstrate the behaviour of this observables obtained in the PWIA and PWIA+FSI with and without the D-wave in the DWF. The significant dependence of the calculation results on the elementary NN-amplitudes is also shown. We conclude with sect. 5.

2 Theoretical formalism

In accordance to the three-body collision theory, let us write the matrix element of the deuteron proton charge exchange reaction,

$$p(\boldsymbol{p}) + d(\boldsymbol{0}) \rightarrow n(\boldsymbol{p}_1) + p(\boldsymbol{p}_2) + p(\boldsymbol{p}_3), \tag{1}$$

in the following form [10]:

$$U_{pd\rightarrow npp} = \sqrt{2}\langle 123|[1-(2,3)] \times[1+t_{23}(E-E_1)g_{23}(E-E_1)]t_{12}^{\mathrm{sym}}|1(23)\rangle, \tag{2}$$

where the operator $g_{23}(E-E_1)$ is a free propagator for the (23)-subsystem and the scattering operator $t_{23}(E-E_1)$ satisfies the Lippmann-Schwinger (LS) equation with the two-body force operator V_{23} as driving term:

$$t_{23}(E-E_1) = V_{23} + V_{23}g_{23}(E-E_1)t_{23}(E-E_1). \tag{3}$$

Here E is the total energy of the three-nucleon system $E = E_1 + E_2 + E_3$.

Let us rewrite the matrix element (2) indicating explicitly the particle quantum numbers,

$$U_{pd\rightarrow npp} = \sqrt{2}\langle \boldsymbol{p}_1 m_1\tau_1, \boldsymbol{p}_2 m_2\tau_2, \boldsymbol{p}_3 m_3\tau_3| \times[1-(2,3)]\omega_{23}t_{12}^{\mathrm{sym}}|\boldsymbol{p}m\tau, \psi_{1M_d00}(23)\rangle, \tag{4}$$

where $\omega_{23} = [1+t_{23}(E-E_1)g_{23}(E-E_1)]$ and the spin and isospin projections are denoted as m and τ, respectively. The permutation operator for two nucleons (i,j) was introduced here. The operator t_{12}^{sym} is the symmetrized NN-operator, $t_{12}^{\mathrm{sym}} = [1-(1,2)]t_{12}$. Inserting the unity

$$\boldsymbol{1} = \int \mathrm{d}\boldsymbol{p}'|\boldsymbol{p}'m'\tau'\rangle\langle\boldsymbol{p}'m'\tau'|,$$

we get the following expression for the reaction amplitude:

$$\begin{aligned}\mathcal{J} = {} & (-1)^{1/2+\tau_3'}\left\langle\frac{1}{2}\tau_2\frac{1}{2}\tau_3\middle|T\tau_2+\tau_3\right\rangle\left\langle\frac{1}{2}\tau_2'\frac{1}{2}\tau_3'\middle|T\tau_2+\tau_3\right\rangle\\ & \times\left\langle\frac{1}{2}\tau_1\frac{1}{2}\tau_2'\middle|T'M_{T'}\right\rangle\left\langle\frac{1}{2}\tau\frac{1}{2}-\tau_3'\middle|T'M_{T'}\right\rangle\\ & \times\left\langle\frac{1}{2}m_2\frac{1}{2}m_3\middle|SM_S\right\rangle\left\langle\frac{1}{2}m_2'\frac{1}{2}m_3'\middle|SM_S'\right\rangle\\ & \times\int\mathrm{d}\boldsymbol{p}_0'\left\langle\boldsymbol{p}_0, SM_S\middle|1+m_N\frac{t^{ST}(E-E_1)}{\boldsymbol{p}_0{}^2-\boldsymbol{p}_0'^2+i0}\middle|\boldsymbol{p}_0', SM_S'\right\rangle\\ & \times\langle\boldsymbol{p}_1 m_1, (\boldsymbol{p}_0'+\boldsymbol{q}/2)\ m_2'|t_{\mathrm{sym}}^{T'}(E-E_3')\\ & \qquad |\boldsymbol{p}m, (\boldsymbol{p}_0'-\boldsymbol{q}/2)\ m''\rangle\\ & \times\langle m''m_3'|\psi_{1M_d}(\boldsymbol{p}_0'-\boldsymbol{q}/2)\rangle\ -\ (2\leftrightarrow 3),\end{aligned} \tag{5}$$

where $E_3' = \sqrt{m_N^2+(\boldsymbol{p}_0'-\boldsymbol{q}/2)^2}$, m_N is the nucleon mass, and we have introduced the momentum transfer $\boldsymbol{q} = \boldsymbol{p}-\boldsymbol{p}_1$, and relative momenta $\boldsymbol{p}_0 = \frac{1}{2}(\boldsymbol{p}_2-\boldsymbol{p}_3)$ and $\boldsymbol{p}_0' = \frac{1}{2}(\boldsymbol{p}_2'-\boldsymbol{p}_3')$. Henceforth, all summations over dummy discrete indices are implied.

In the momentum representation the DWF $\psi_{1M_d}(\boldsymbol{k})$ with spin projection M_d is written as

$$|\psi_{1M_d}(\boldsymbol{k})\rangle = \sum_{L=0,2}\sum_{M_L=-L}^{L}\langle LM_L1M_s|1M_d\rangle \times u_L(k)Y_L^{M_L}(\hat{k})|1M_s\rangle, \tag{6}$$

with the spherical harmonics $Y_L^{M_L}(\hat{k})$ and the Clebsh-Gordon coefficients in standard form. In our calculations, we have employed the following parameterizations of the S- and D-state wave functions:

$$u_0(p) = \sqrt{\frac{2}{\pi}}\sum_i\frac{C_i}{\alpha_i^2+p^2}\,, \quad u_2(p) = \sqrt{\frac{2}{\pi}}\sum_i\frac{D_i}{\beta_i^2+p^2} \tag{7}$$

as proposed in refs. [13–15] .

We assume that $\tau = \tau_2 = \tau_3 = 1/2$ and $\tau_1 = -1/2$. Then the isotopic coefficient can be calculated and eq. (5) is simplified as

$$\begin{aligned}\mathcal{J} = {} & \frac{1}{2}\left\langle\frac{1}{2}m_2\frac{1}{2}m_3\middle|SM_S\right\rangle\left\langle\frac{1}{2}m_2'\frac{1}{2}m_3'\middle|SM_S'\right\rangle\\ & \times\left\langle LM_L1\mathcal{M}_{\mathcal{S}}\middle|1M_d\right\rangle\left\langle\frac{1}{2}m''\frac{1}{2}m_3'\middle|1\mathcal{M}_{\mathcal{S}}\right\rangle\\ & \times\int\mathrm{d}\boldsymbol{p}_0'\left\langle\psi_{\boldsymbol{p}_0SM_STM_T}^{(-)}\middle|\boldsymbol{p}_0'SM_S'TM_T\right\rangle\\ & \times u_L(|\boldsymbol{p}_0'-\boldsymbol{q}/2|)Y_L^{M_L}(\widehat{\boldsymbol{p}_0'-\boldsymbol{q}}/2)\\ & \times\langle\boldsymbol{p}_1 m_1, (\boldsymbol{p}_0'+\boldsymbol{q}/2)\ m_2'|t_{\mathrm{sym}}^0(E-E_3')\\ & -t_{\mathrm{sym}}^1(E-E_3')|\boldsymbol{p}m, (\boldsymbol{p}_0'-\boldsymbol{q}/2)\ m''\rangle\ -\ (2\leftrightarrow 3).\end{aligned} \tag{8}$$

The wave function of the final pp-pair,

$$\left\langle\psi_{\boldsymbol{p}_0SM_STM_T}^{(-)}\middle|\boldsymbol{p}_0'SM_S'TM_T\right\rangle = \delta(\boldsymbol{p}_0-\boldsymbol{p}_0')\delta_{M_sM_s'} + \frac{m_N}{\boldsymbol{p}_0^2-\boldsymbol{p}_0'^2+i0}\langle\boldsymbol{p}_0SM_S|t^{ST}|\boldsymbol{p}_0'SM_S'\rangle\,, \tag{9}$$

contains the FSI part, which can be taken in different ways.

In this paper we use the matrix inversion method (MIM) suggested in refs. [11,12] and applied to study the deuteron electro-disintegration [16,17] and the deuteron proton breakup process [10]. As in ref. [16], we consider the truncated partial-wave expansion,

$$\begin{aligned}\left\langle\psi_{\boldsymbol{p}_0SM_STM_T}^{(-)}\middle|\boldsymbol{p}_0'SM_S'TM_T\right\rangle = {} & \delta_{M_SM_S'}\delta(\boldsymbol{p}_0-\boldsymbol{p}_0')\\ & +\sum_{J=0}^{J_{\max}}\sum_{M_J=-J}^{J}Y_l^{\mu}(\hat{p}_0)\langle l\mu SM_S|JM_J\rangle\psi_{ll'}^{\alpha}(p_0')\\ & \times\langle l'\mu'SM_S'|JM_J\rangle Y_{l'}^{*\mu'}(\hat{p}_0'),\end{aligned} \tag{10}$$

where $J_{\max}$ is the maximum value of the total angular momentum in the pp-partial waves and $\alpha = \{J, S, T\}$ is the set of conserved quantum numbers. The radial functions $\psi^{\alpha}_{ll'}(p'_0)$ are related via

$$\psi^{\alpha}_{ll'}(p'_0) = \sum_{l''} O_{ll''} \varphi^{\alpha}_{l''l'}(p'_0) - \frac{\delta(p'_0 - p_0)}{p_0^2} \delta_{ll'} \tag{11}$$

to the partial-wave functions $\varphi^{\alpha}_{l''l'}(p'_0)$, which have the asymptotics of standing waves. The coefficients $O_{ll''}$ can be expressed in terms of the corresponding phase shifts and mixing parameters [16].

Within the MIM, the functions $\varphi^{\alpha}_{ll'}(p'_0)$ can be represented as

$$\varphi^{\alpha}_{ll'}(p'_0) = \sum_{j=1}^{N+1} B^{\alpha}_{ll'}(j) \frac{\delta(p'_0 - p_j)}{p_j^2}, \tag{12}$$

where the coefficients $B^{\alpha}_{ll'}(j)$ fulfill a set of linear algebraic equations approximately equivalent to the LS integral equation for the pp scattering problem[1]. Here N is the dimension of this set, p_j are the grid points associated with the Gaussian nodes over the interval $[-1, 1]$ and $p_{N+1} = p_0$ (details can be found in ref. [18]). It should be noted that in this way the nucleon wave function is expressed by a series of δ-functions allowing one to reduce a triple integral in eq. (8) to a double one. In addition, the method offers the opportunity to consider the nucleon wave function in the continuum directly in momentum space which simplifies all subsequent calculations.

3 Collinear geometry

In this paper we consider the special kinematics, when the momentum transfer $\boldsymbol{q} = \boldsymbol{p} - \boldsymbol{p}_1$ is close to zero. In other words, the neutron momentum has the same value and direction as the beam proton. In fact, the momentum transfer is not exactly zero, $q \approx 1.8$ MeV$/c$ due to the difference between proton and neutron masses and deuteron binding energy. But since this value is very small and has no significant influence on the results, we shall suppose $q = 0$ in the subsequent calculations.

Under such kinematical conditions one can neglect the dependence of the high-energy nucleon-nucleon matrix $t_{NN}(E - E'_3)$ in eq. (8) on the internal nucleon-nucleon momentum in the deuteron and express it in the center-of-mass system (c.m.s.) through three independent amplitudes:

$$t^{\mathrm{cm}}_{NN}(\boldsymbol{q} = 0) = A + (F - B)(\boldsymbol{\sigma}_1 \hat{Q}^*)(\boldsymbol{\sigma}_2 \hat{Q}^*) + B(\boldsymbol{\sigma}_1 \boldsymbol{\sigma}_2), \tag{13}$$

where

$$\hat{Q}^* = \frac{\boldsymbol{p}^* + \boldsymbol{p}'^*}{|\boldsymbol{p}^* + \boldsymbol{p}'^*|} = \hat{p}^*$$

with fast-proton (neutron) momentum $\boldsymbol{p}^*$. We use some results of the relativistic potential theory [19,20] to relate this NN t-matrix in the c.m.s. with that in the frame of interest (see, also [10]):

$$\begin{aligned} \langle m_1 m'_2, \boldsymbol{p}\boldsymbol{p}'_0 | t | \boldsymbol{p}\boldsymbol{p}'_0, mm'' \rangle &= NN'F \\ &\times \langle m_1 | D^\dagger(\boldsymbol{u}, \boldsymbol{p}) | \mu_1 \rangle \langle m'_2 | D^\dagger(\boldsymbol{u}, \boldsymbol{p}'_0) | \mu'_2 \rangle \\ &\times \langle \mu_1 \mu'_2 | t_{\mathrm{cm}}(\boldsymbol{p}^*) | \mu \mu'' \rangle \langle \mu | D(\boldsymbol{u}, \boldsymbol{p}) | m \rangle \\ &\times \langle \mu'' | D(\boldsymbol{u}, \boldsymbol{p}'_0) | m'' \rangle, \end{aligned} \tag{14}$$

where D is the Wigner rotation operator in the spin space and u is the four-velocity. In our kinematical situation, when $\boldsymbol{p} = \boldsymbol{p}_1 \gg \boldsymbol{p}'_0$, each of these operators is slightly different from the unit operator, so that with a good approximation the t_{NN}-matrix in the frame of interest has the same spin structure. The product of the normalization factors N and N' and the kinematical factor F is

$$NN'F = \frac{m_N + E_p}{4E_p}. \tag{15}$$

In such a way, we have the following relation for the high-energy NN t-matrix in the different frames of reference:

$$\begin{aligned} \langle m_1 m'_2 | t(\boldsymbol{p}, \boldsymbol{p}'_0) | mm'' \rangle &= \frac{m_N + E_p}{4E_p} \\ &\times \langle m_1 m'_2 | t_{\mathrm{cm}}(\boldsymbol{p}^*) | mm'' \rangle. \end{aligned} \tag{16}$$

To evaluate such quantities without their momentum angular decomposition, we use the phenomenological model suggested by Love and Franey in refs. [21]. In this approach the corresponding matrix elements are expressed through the effective NN-interaction operators sandwiched between the initial and final plane-wave states, that enables us to extend this construction to the off-shell case. Obviously, such off-shell extrapolation does not change the general spin structure.

Since the pp-pair belongs to the isotriplet, one can anticipate that the FSI in the ${}^1S_0$-state is prevalent at comparatively small p_0 values. In such a way we get the following expression for amplitude of the dp charge exchange process:

$$\begin{aligned} \mathcal{J} &= \mathcal{J}_{\mathrm{PWIA}} + \mathcal{J}_{{}^1S_0} \\ \mathcal{J}_{\mathrm{PWIA}} &= \frac{m_N + E_p}{4E_p} \langle LM_L 1\mathcal{M}_{\mathcal{S}} | 1M_D \rangle u_L(p_0) Y_L^{M_L}(\widehat{p}_0) \\ &\times \Big\{ \Big\langle \frac{1}{2} m'_2 \frac{1}{2} m_3 | 1\mathcal{M}_{\mathcal{S}} \Big\rangle \langle m_1 m_2 | t^0_{\mathrm{cm}}(\boldsymbol{p}^*) - t^1_{\mathrm{cm}}(\boldsymbol{p}^*) | mm'_2 \rangle \\ &- \Big\langle \frac{1}{2} m'_2 \frac{1}{2} m_2 | 1\mathcal{M}_{\mathcal{S}} \Big\rangle \langle m_1 m_3 | t^0_{\mathrm{cm}}(\boldsymbol{p}^*) - t^1_{\mathrm{cm}}(\boldsymbol{p}^*) | mm'_2 \rangle \Big\}, \\ \mathcal{J}_{{}^1S_0} &= \frac{(-1)^{1-m_2-m'_2}}{4\pi} \frac{m_N + E_p}{4E_p} \delta_{m_2\ -m_3} \\ &\times \langle LM_L 1\mathcal{M}_{\mathcal{S}} | 1M_D \rangle \Big\langle \frac{1}{2} m'' \frac{1}{2} - m'_2 | 1\mathcal{M}_{\mathcal{S}} \Big\rangle \\ &\times \int \mathrm{d}p'_0 p'^2_0 \int \mathrm{d}\hat{p}'_0 u_L(p'_0) Y_L^{M_L}(\widehat{p'_0}) \psi^{001}_{00}(p'_0) \\ &\times \langle m_1 m'_2 | t^0_{\mathrm{cm}}(\boldsymbol{p}^*) - t^1_{\mathrm{cm}}(\boldsymbol{p}^*) | mm'' \rangle. \end{aligned} \tag{17}$$

As one can see, we have a very simple integral over the angular variable $\hat{p}'_0$. As a result of this integration we get

[1] We neglect here the Coulomb interaction in the pp-pair.

the following relation for $\mathcal{J}_{^1S_0}$:

$$\mathcal{J}_{^1S_0} = \frac{(-1)^{1-m_2-m_2'}}{\sqrt{4\pi}} \frac{m_N + E_p}{4E_p} \delta_{m_2-m_3} \times \left\langle \frac{1}{2}m'' \frac{1}{2} - m_2' | 1 M_D \right\rangle \langle m_1 m_2' | t^0_{\rm cm}(\boldsymbol{p}^*) - t^1_{\rm cm}(\boldsymbol{p}^*) | m m'' \rangle \times \int {\rm d}p_0' p_0'^2 \psi^{001}_{00}(p_0') u_0(p_0') \,. \quad (18)$$

Note that the integral over the radial variable p_0' also does not present any difficulties, since $\psi^{001}_{00}(p_0')$ contains δ-functions.

4 Results and discussions

We define the unpolarized $2 \to 3$ cross-section by the standard manner

$$\sigma(dp \to npp) = (2\pi)^4 \frac{E_p}{p} \int {\rm d}\boldsymbol{p}_1 {\rm d}\boldsymbol{p}_2 {\rm d}\boldsymbol{p}_3 |\overline{\mathcal{J}}|^2 \times \delta^4(\text{4-momentum}), \quad (19)$$

where $|\overline{\mathcal{J}}|^2 = 1/6{\rm Tr}(\mathcal{J}\mathcal{J}^+)$ is the square of the process amplitude averaged over all particles spin states. Using the δ^3-function to eliminate the $\boldsymbol{p}_3$ integration and changing variables from $\boldsymbol{p}_1$ to $\boldsymbol{q}$, we have

$$\sigma(dp \to npp) = (2\pi)^4 \frac{E_p}{p} \int {\rm d}\boldsymbol{q} {\rm d}\boldsymbol{p}_2 |\overline{\mathcal{J}}|^2 \delta\Bigg(m_d + E_p - \sqrt{m_N^2 + (\boldsymbol{p}-\boldsymbol{q})^2} - E_2 - \sqrt{m_N^2 + (\boldsymbol{q}-\boldsymbol{p}_2)^2} \Bigg) . \quad (20)$$

Taking $p_2, q \ll p$, this expression can be reduced to

$$\sigma(dp \to npp) = (2\pi)^6 \frac{E_p^2}{2p^2} \int {\rm d}q^2 {\rm d}p_2 {\rm d}\cos\theta_2 p_2^2 |\overline{\mathcal{J}}|^2 . \quad (21)$$

We define the general spin observable related to the polarization of initial particles in terms of the Pauli 2×2 spin matrices σ for the proton and a set of spin operators S for the deuteron [22] as follows:

$$C_{\alpha\beta} = \frac{{\rm Tr}(\mathcal{J}\sigma_\alpha S_\beta \mathcal{J})}{{\rm Tr}(\mathcal{J}\mathcal{J}^+)} , \quad (22)$$

where the indices α and β refer to the proton and deuteron polarization, respectively; σ_0 and S_0 corresponding to the unpolarized particles are the unit matrices of two and three dimensions. In such a way, eqs. (17)-(18) for the dp charge exchange amplitude enable us to get the relation for any variable of this process taking into account two-slow-protons final-state interaction in the $^1S_0$-state. So, we have the following expression for the spin-averaged squared amplitude:

$$C_0 \equiv {\rm Tr}(\mathcal{J}\mathcal{J}^+) = \frac{1}{4\pi} \left(\frac{m_N + E_p}{2E_p} \right)^2 \Big\{ 2(2B^2 + F^2) \times (\mathcal{U}^2(p_2) + w^2(p_2)) + (F^2 - B^2) w(p_2) \times (w(p_2) - 2\sqrt{2}{\rm Re}\,\mathcal{U}(p_2))(3\cos^2\theta_2 - 1) \Big\} , \quad (23)$$

where $\mathcal{U}(p_2) = u(p_2) + \int {\rm d}p_0' {p_0'}^2 \psi^{001}_{00}(p_0') u(p_0')$ is the S-component of the DWF corrected on the FSI of the pp-pair. As one can see, C_0 is independent of the neutron and proton azimuthal angles ϕ_q and ϕ_2. To obtain eq. (21), we have considered this fact and performed the integration over azimuthal angles. We use a right-hand coordinate system defined in accordance to the Madison convention [23]. The quantization z-axis is along the beam proton momentum $\boldsymbol{p}$. Since the direction of $\boldsymbol{p} \times \boldsymbol{p}_1$ is undefined in the collinear geometry, we choose the y-axis normal to the beam momentum. Then the third axis is $\boldsymbol{x} = \boldsymbol{y} \times \boldsymbol{z}$.

The tensor analyzing power can be presented in the following form:

$$C_{0,yy} \cdot C_0 = \frac{1}{2} \frac{1}{4\pi} \left(\frac{m_N + E_p}{2E_p} \right)^2 \times \Big\{ 4(F^2 - B^2)(\mathcal{U}^2(p_2) + w^2(p_2)) + (2F^2 + B^2) w(p_2) \times (w(p_2) - 2\sqrt{2}{\rm Re}\,\mathcal{U}(p_2))(3\cos^2\theta_2 - 1) + 9B^2 w(p_2)(w(p_2) - 2\sqrt{2}{\rm Re}\,\mathcal{U}(p_2)) \sin^2\theta_2 \cos 2\phi_2 - 54(F^2 - B^2) w^2(p_2) \sin^2\theta_2 \cos^2\theta_2 \sin^2\phi_2 \Big\} . \quad (24)$$

Note that only the squared nucleon-nucleon spin-flip amplitudes B^2 and F^2 are in the expression for the tensor analyzing power $C_{0,yy}$ and the differential cross-section. However, the spin correlation due to the vector polarization of the deuteron and beam proton contains the interference terms of these amplitudes:

$$C_{y,y} \cdot C_0 = -\frac{2}{4\pi} \left(\frac{m_N + E_p}{2E_p} \right)^2 \times \{ {\rm Re}(FB^*) \{ 2\mathcal{U}^2(p_2) - w^2(p_2) - w(p_2) \times (w(p_2) + \sqrt{2}{\rm Re}\,\mathcal{U}(p_2))(1 - 3\sin^2\theta_2 \sin^2\phi_2) \} + 12 w(p_2) {\rm Im}(FB^*) {\rm Im}\,\mathcal{U}(p_2) \times (\cos^2\theta_2 - \sin^2\theta_2 \cos^2\phi_2) \} . \quad (25)$$

It is interesting that there is a term proportional to the imaginary part of $\mathcal{U}(p_2)$. It has a non-zero value only in the case when the FSI is taken into account. An analogous result we have obtained for the vector-tensor spin correlation:

$$C_{y,xz} \cdot C_0 = -\frac{3}{4\pi} \left(\frac{m_N + E_p}{2E_p} \right)^2 \times \{ {\rm Im}(FB^*) \{ 2\mathcal{U}^2(p_2) - w^2(p_2) - w(p_2) \times (w(p_2) + \sqrt{2}{\rm Re}\,\mathcal{U}(p_2))(1 - 3\sin^2\theta_2 \sin^2\phi_2) + 18 w^2(p_2) \sin^2\theta_2 \cos^2\theta_2 \cos^2\phi_2 \} - 3\sqrt{2} w(p_2) {\rm Re}(FB^*) {\rm Im}\,\mathcal{U}(p_2) \times (\cos^2\theta_2 - \sin^2\theta_2 \cos^2\phi_2) \} . \quad (26)$$

In order to evaluate these observables we consider the kinematics, when one of the slow protons is emitted along the beam direction as well as the neutron, *i.e.* $\theta_2 = 0°$.

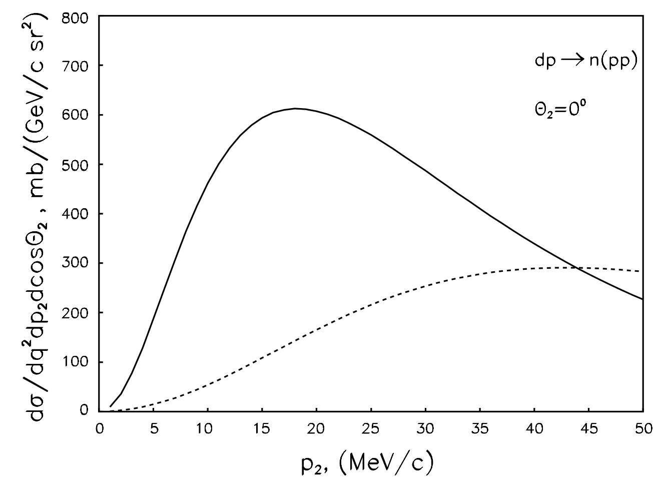

Fig. 1. The differential cross-section at $\boldsymbol{q} = 0$ as a function of one of the slow-proton momentum. The dashed and full lines correspond to the PWIA and PWIA+FSI, respectively.

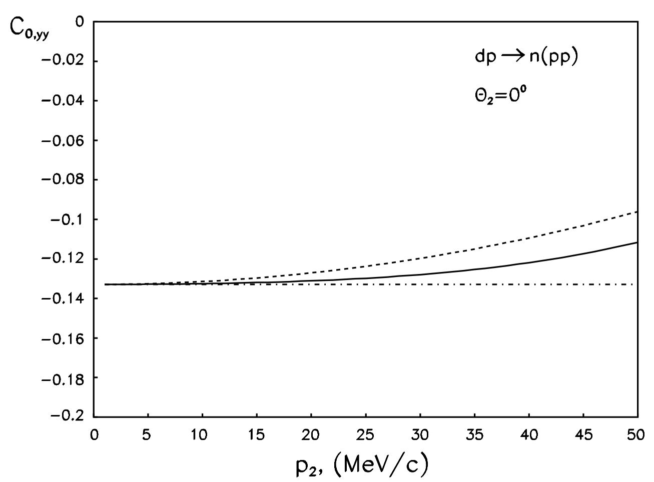

Fig. 2. The tensor analyzing power C_{yy} *vs.* p_2. The dashed line corresponds to PWIA; dash-dotted and full lines are PWIA+FSI without and with D-component in the DWF, respectively.

Under such conditions eqs. (23)-(26) are significantly simplified:

$$C_0 = \frac{1}{2\pi}\left(\frac{m_N + E_p}{2E_p}\right)^2 \times\{(2B^2+F^2)(\mathcal{U}^2(p_2)+w^2(p_2)) +(F^2-B^2)w(p_2)(w(p_2)-2\sqrt{2}\mathrm{Re}\,\mathcal{U}(p_2))\},$$

$$C_{0,yy}\cdot C_0 = \frac{1}{4\pi}\left(\frac{m_N + E_p}{2E_p}\right)^2 \times\{2(F^2-B^2)(\mathcal{U}^2(p_2)+w^2(p_2)) +(2F^2+B^2)w(p_2)(w(p_2)-2\sqrt{2}\mathrm{Re}\,\mathcal{U}(p_2))\},$$

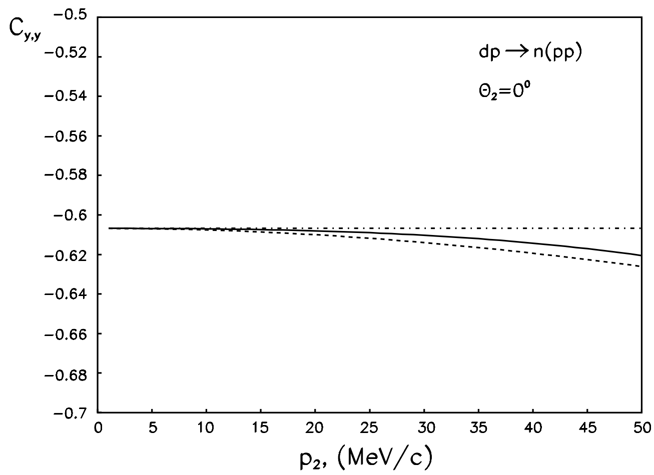

Fig. 3. The spin-correlation $C_{y,y}$ due to the vector polarization of the deuteron. The curves are the same as in fig. 2.

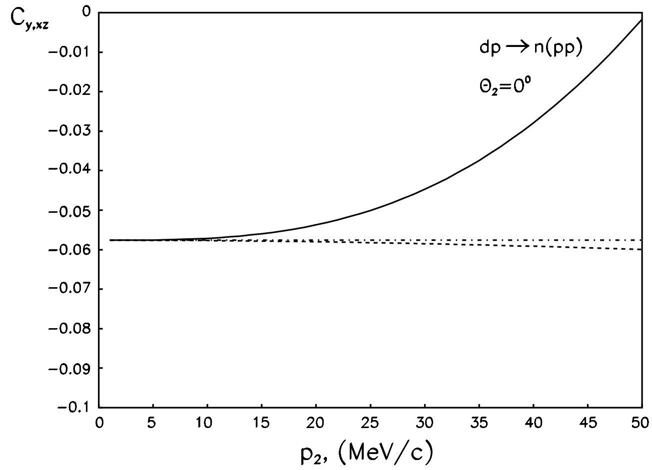

Fig. 4. The spin-correlation $C_{y,xz}$ due to the tensor polarization of the deuteron. The curves are the same as in fig. 2.

$$C_{y,y}\cdot C_0 = -\frac{2}{4\pi}\left(\frac{m_N + E_p}{2E_p}\right)^2 \Big\{\mathrm{Re}(FB^*)[2\mathcal{U}^2(p_2) -2w^2(p_2)-\sqrt{2}\mathrm{Re}\,\mathcal{U}(p_2)w(p_2)] -3\sqrt{2}\mathrm{Im}(FB^*)\mathrm{Im}\,\mathcal{U}(p_2)w(p_2)\Big\},$$

$$C_{y,xz}\cdot C_0 = -\frac{3}{4\pi}\left(\frac{m_N + E_p}{2E_p}\right)^2 \Big\{\mathrm{Im}(FB^*)[2\mathcal{U}^2(p_2) -2w^2(p_2)-\sqrt{2}\mathrm{Re}\,\mathcal{U}(p_2)w(p_2)] +3\sqrt{2}\mathrm{Re}(FB^*)\mathrm{Im}\,\mathcal{U}(p_2)w(p_2)\Big\}. \quad (27)$$

The differential cross-section and three polarization observables are presented in figs. 1-4. The Love and Franey parametrization with a set of parameters obtained by fitting of the modern phase shift data SP00 [25,26] has been employed for the NN-amplitude. The full lines correspond to calculations taking into account both the FSI in the *pp*-pair and S- and D-waves in the deuteron. The results obtained in the PWIA are shown by the dashed lines.

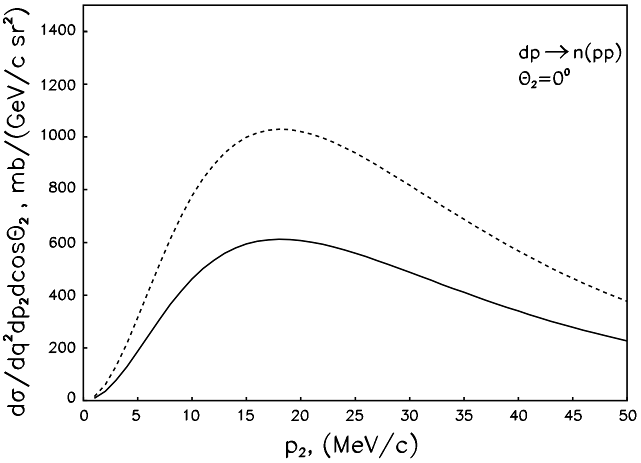

Fig. 5. The differential cross-section at $\boldsymbol{q} = 0$ as a function of p_2. The dashed and full lines correspond to the full calculation with a set of NN amplitude parameters taken from [21] and a fit of SP00 data [26].

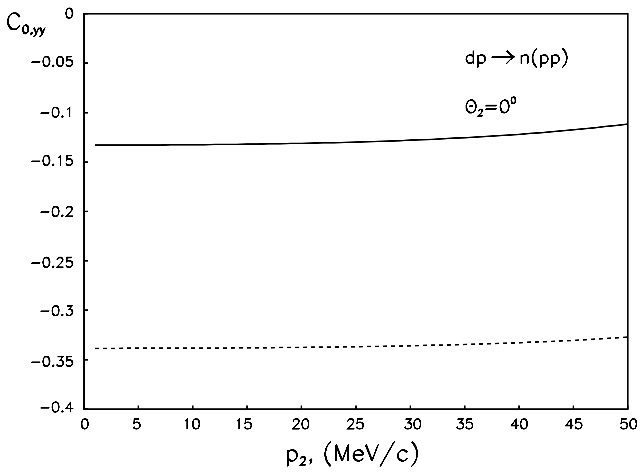

Fig. 6. The tensor analyzing power C_{yy} *vs.* p_2. The curves are the same as in fig. 5.

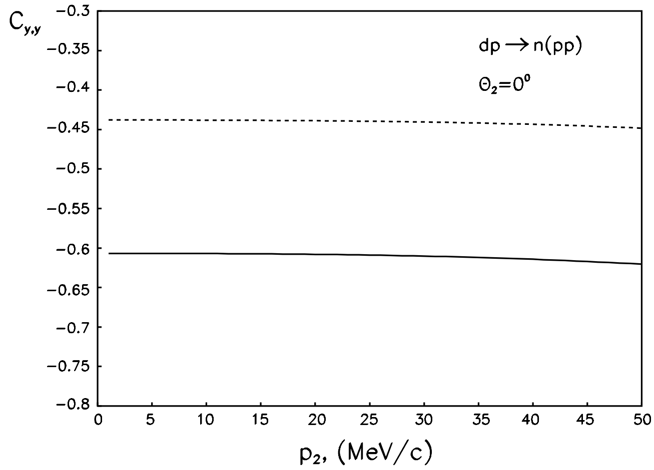

Fig. 7. The spin-correlation $C_{y,y}$ due to the vector polarization of the deuteron. The curves are the same as in fig. 5.

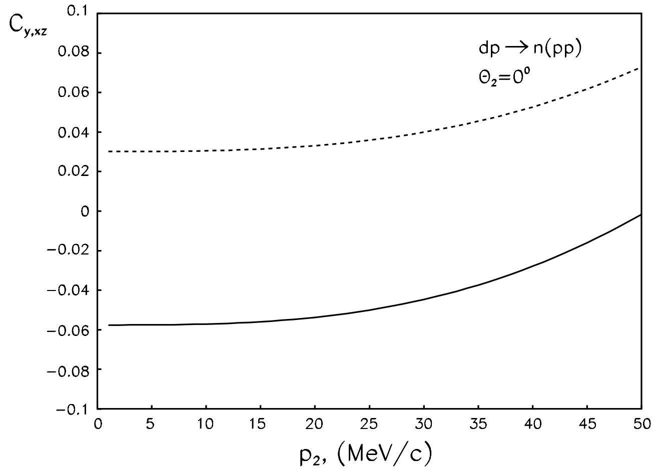

Fig. 8. The spin-correlation $C_{y,xz}$ due to the tensor polarization of the deuteron. The curves are the same as in fig. 5.

The dash-dotted lines in figs. 2-4 for polarization observables are the full calculation results without D-wave in the DWF. In fig. 1 the full and dash-dotted lines are undistinguished. All calculations were carried out with Paris NN-potential [24] and Paris DWF [13].

As one can see, the FSI contribution to the differential cross-section (fig. 1) is significant even at very small proton momentum, while for the polarization observables the difference between PWIA and PWIA+FSI is visible only for $p_2 \geq 10$–15 MeV/c. However, with the increase of the proton momentum up to 50 MeV/c the importance of the FSI corrections to the PWIA also increases.

Note that the absolute value of the tensor analyzing power $C_{0,yy}$ (fig. 2) in the momentum interval of interest is near zero. In order to understand the source of that, we disregard the D-wave in the DWF. Then, as a consequence from eq. (27), the polarization observables are defined by the ratio of the nucleon-nucleon charge exchange amplitudes only

$$\begin{aligned} C_{0,yy} &= \frac{1}{2} \cdot \frac{F^2 - B^2}{2B^2 + F^2}\,, \\ C_{y,y} &= -2 \cdot \frac{\mathrm{Re}(FB^*)}{2B^2 + F^2}\,, \\ C_{y,xz} &= -3 \cdot \frac{\mathrm{Im}(FB^*)}{2B^2 + F^2}\,. \end{aligned} \tag{28}$$

Thus, the nearness of the tensor analyzing power to zero indicates that the absolute values of the spin-flip NN amplitudes approximately equal each other, $|B| \approx |F|$.

The vector-tensor spin correlation $C_{y,xz}$ (fig. 4) has also very small value, $|C_{y,xz}| \approx 0.06$. The magnitude of this observable decreases down to zero for $p_2 \approx 50$ MeV/c, if the FSI corrections and D-wave in the deuteron are taken into account, while it is almost constant in the PWIA and PWIA+FSI without D-wave. As one can see from eqs. (27), (28) for $C_{y,xz}$, the reason of this behaviour is connected with the small value of the imaginary part of the nucleon-nucleon amplitudes product, $\mathrm{Im}(FB^*)$. In such a way, the great contribution to $C_{y,xz}$ gives a term proportional to $\mathrm{Re}(FB^*)$, which is defined by the D-wave and the imaginary part of the generalized function $\mathcal{U}(p_2)$. Note that $\mathrm{Im}\,\mathcal{U}(p_2) \neq 0$, if FSI taken into account.

The other situation is for the vector-vector spin correlation $C_{y,y}$ (fig. 3). The term proportional to $\mathrm{Re}(FB^*)$ gives also a considerable contribution to this observable, but it is multiplied by the $\mathcal{U}^2(p_2)$. The magnitude of $C_{y,y}$ is close to the theoretical limit $-2/3$, that confirms to the conclusion about the approximate equality of the nucleon-nucleon amplitudes, $|B|$ and $|F|$. Besides, this allows to conclude, that the relative phase between these amplitudes is close to zero. This is seen from eq. (28), where the D-wave was neglected.

Since all the considered observables are defined by the elementary nucleon-nucleon amplitudes mostly, it is interesting to compare their behaviour for different NN parametrization. In figs. 5-8 we present the same observables as in figs. 1-4 for two sets of parameters. The full line corresponds to the parameterization based on the modern shift analysis SP00 [25,26]. The dashed line is obtained using a set of parameters for the NN amplitude from [21]. As one can see, the difference for the differential cross-section (fig. 5) is about 1.5–2 times. The absolute value for the tensor analyzing power $C_{0,yy}$ (fig. 6) with the new parametrization is about 2.5 times smaller than that with parametrization [21]. The opposite situation is for the vector-vector spin correlation $C_{y,y}$ (fig. 7), where the new prediction is 1.5 times lager in comparison with the old parametrization of the NN amplitude [21]. The predictions of these two parameterizations for the vector-tensor spin correlation $C_{y,xz}$ (fig. 8) are even opposite in sign. Nevertheless, the qualitative behaviour of the curves in figs. 6-8 for different sets of parameters is similar.

5 Conclusion

We have studied the deuteron-proton charge exchange reaction at 1 GeV energy in special kinematics, $\boldsymbol{q} \approx 0$. The influence of the D-wave in the deuteron and FSI between two slow protons has been considered. It was shown, that D-wave and FSI effects are negligible for the polarization observables at proton momentum up to 10–15 MeV/c. As a result, in this region the polarization observables are defined by the ratio of the nucleon-nucleon charge exchange amplitudes only. However, it should not be ignored that the importance of the D-wave and, especially, FSI in polarization observables increases at $p_2 \geq 15$ MeV/c. In such a way, we conclude that the ratio of the nucleon-nucleon charge exchange amplitudes and the phase shift between them can be extracted from experimental data rather simply, if the experimental conditions and technical setup possibilities allow to work in this small momentum interval. In the opposite case, this procedure is more complicated and model dependent. It should be remembered that the FSI contribution to the differential cross-section is very significant in comparison with PWIA predictions even at very small proton momentum. This fact does not enable us to get the absolute value of the nucleon-nucleon spin flip amplitudes without considering the FSI corrections.

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