Threshold open charm production in nucleon-nucleon collisions

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Abstract. The associative charm particles production in nucleon-nucleon collisions $N + N \to \Lambda_c(\Sigma_c) + \overline{D} + N$, is described in a general way and the spin and isospin structure of the corresponding matrix elements are derived. Using an analogy with strange-particle (ΛK) production, the D-meson exchange model is considered in detail. Estimations of the energy behavior of the threshold cross-sections show a large dependence on the form of the phenomenological hadron form factors and indicate that, at threshold, the cross-section is three orders of magnitude smaller than for strange-particle production.

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1 Introduction

New experimental facilities, planned in the near future [1], will be able to make a detailed measurement of open charm production in pN, NA or AA collisions, in a wide energy region, starting from threshold. One can expect that the physics of close-to-threshold $\Lambda_c D$ production is similar to ΛK production.

The last precise experiments about strange-particle production in pp collisions gave interesting information about the possible reaction mechanism [2]. For example, the relative value of the total cross-sections for ΛK and $\Sigma^0 K$ production, and in particular the large and negative polarization transfer coefficient D_{nn} (from the polarized proton beam to the produced Λ-hyperon) can be considered a strong indication of the K-meson exchange as the main reaction mechanism —in this energy region. The final ΛN interaction strongly affects the different polarization observables, suggesting a sensitive method for the determination of low-energy parameters of the ΛN or ΣN interaction. Such information is useful for the determination of the ΛN and ΣN potential and therefore for predictions of the properties of hypernuclei [3].

The existing data about $N + N \to \Lambda(\Sigma) + K + N$, in the near-threshold region, do not exclude a contribution of the π exchange, in particular in the case of Σ production, where the K exchange is depressed by a smaller $g_{p\Sigma K}$

coupling constant (following the SU_3 symmetry one finds $g^2_{\Lambda pK}/g^2_{\Sigma pK} \simeq 27$).

The experience gained in the theoretical and experimental study of strange particles can be extrapolated to the case of the processes $N + N \to \Lambda_c(\Sigma_c) + \overline{D} + N$, in the threshold region. No experimental data exist here: the lowest energy, where open charm production has been detected in pp collisions is $E_p = 70$ GeV ($\sqrt{s} = 11$ GeV) [4]. The theoretical activity is focused on QCD-inspired approaches, mainly for inclusive charm production, which cannot be easily extrapolated to the threshold region. In a more traditional way, the process $p + p \to \Lambda_c + \overline{D} + N$ has been considered twenty years ago, in the framework of one-pion exchange model [5], using a specific approach, for the calculation of the amplitudes for such elementary subprocess, as $\pi + N \to \Lambda_c + D$ or $\pi + N \to \Lambda_c + D^*$. This approach is based on the Quark-Gluon-String model, which can be considered as the non-perturbative topological equivalent of the corresponding Regge phenomenology. Note that, in such approach, the spin structure of the amplitudes for $\pi + N \to \Lambda_c + D(D^*)$ process is ignored, and in principle it is crucial for a correct calculation in the near-threshold region. Moreover, a number of parameters is necessary for such calculations, concerning, in particular, the D^* Regge trajectory.

The understanding of elementary open charm production processes in πN or NN collisions is a very actual problem, in connection with leptonic pair production in Pb + Pb collisions [6] at SPS. The intensive charm production in secondary πN collisions [7] can be considered

an important source of leptonic pairs in the intermediate-mass region [8].

For a reliable description of the overall charm dynamics of AA collisions, it is necessary to have a good parametrization of the elementary processes of open charm production in πN and NN collisions. The parametrization of the energy behavior of the total cross-section for different processes of D and $\overline{D}$ production for $\sqrt{s} \geq 10$ GeV generally used in simulation codes, cannot be applied near threshold as it violates the general threshold behavior for two- or three-particle production. We completely agree on the explicit statement of ref. [9]: *We have to point out that our parametrizations for the differential and total cross-sections for open charm (as well as charmonia) become questionable at low energy, but also at high energy. It is thus mandatory that they have to be controlled by experimental data from pp, pA and πN reactions before reliable conclusions on open charm dynamics in nucleus-nucleus reactions can be drawn.*

In this paper we analyze the $N + N \rightarrow \Lambda_c(\Sigma_c) + \overline{D} + N$ processes, in the near-threshold region, taking rigorously into account the spin and isospin structure of the corresponding matrix elements. In our analysis, we follow a very close analogy between open charm and open strangeness production in nucleon-nucleon collisions. According to the SU_4 symmetry (which gives useful relations for meson-baryon coupling constants, such as $g_{p\Lambda K^+} = g_{p\Lambda_c^+ \overline{D^0}}$ and $g_{p\Sigma^0 K^+} = g_{p\Sigma_c^0 \overline{D^0}}$), one can expect that namely the D exchange is important for the processes $N + N \rightarrow \Lambda_c(\Sigma_c) + \overline{D} + N$ in the near-threshold region. The amplitude for such mechanism can be predicted, being driven by elastic $\overline{D}N$ scattering (at zero energy), in analogy to elastic KN scattering, with real scattering lengths. The possible π exchange —involving the elementary subprocess $\pi + N \rightarrow \Lambda_c + D$, seems more complicated and difficult to estimate (in a model-independent way).

This paper is organized as follows. In sect. 2 we establish the spin and isospin structure of the threshold matrix elements for the processes $N + N \rightarrow \Lambda_c(\Sigma_c) + \overline{D} + N$. The calculation of the threshold partial amplitudes for the $\Lambda_c \overline{D}$ production in the framework of the $\overline{D}$ exchange is done in sect. 3. The estimation of the energy dependence of the total cross-section is done in sect. 4.

2 Spin and isospin in the processes $N + N \rightarrow \Lambda_c(\Sigma_c) + \overline{D} + N$

The strong correlation between the spin and isospin structures of the threshold matrix element for the $N + N \rightarrow \Lambda_c(\Sigma_c) + \overline{D} + N$ reactions is an important characteristic of the threshold NN dynamics. First of all, this correlation causes the absence of factorization of spin and isospin variables, which is typical for the majority of hadronic processes (see πN scattering, for example). Factorization means that we can consider the spin structure of the matrix elements independently of the isospin structure. Moreover, all possible spin structures, which are allowed by the

symmetry properties of the strong interaction for the considered matrix element, contain, in general, all possible isospin structures. It is not the case for the threshold matrix element of the $N + N \rightarrow \Lambda_c(\Sigma_c) + \overline{D} + N$ processes, where there is a strong and not trivial connection of the spin and isospin variables. As a result, different spin amplitudes have different isotopic structures.

Using the isotopic invariance of the strong interaction and taking into account that $I(\Sigma_c) = 1$, $I(\overline{D}) = 1/2$, one can find for $N + N \rightarrow \Sigma_c + \overline{D} + N$:

$$
\begin{aligned}
\mathcal{M}(pp \rightarrow \Sigma_c^{++} D^- p) &= & A_{11} &+ \sqrt{2} A_{10}, \\
\mathcal{M}(pp \rightarrow \Sigma_c^{++} \overline{D^0} n) &= & A_{11} &- \sqrt{2} A_{10}, \\
\mathcal{M}(pp \rightarrow \Sigma_c^{+} \overline{D^0} p) &= -\sqrt{2} A_{11}, \\
\mathcal{M}(np \rightarrow \Sigma_c^{++} D^- n) &= & A_{11} & &+ \sqrt{2} A_{01}, \\
\mathcal{M}(np \rightarrow \Sigma_c^{0} \overline{D^0} p) &= & -A_{11} & &+ \sqrt{2} A_{01}, \\
\mathcal{M}(np \rightarrow \Sigma_c^{+} D^- p) &= & & A_{10} &- A_{01}, \\
\mathcal{M}(np \rightarrow \Sigma_c^{+} \overline{D^0} n) &= & & -A_{10} &- A_{01},
\end{aligned}
\tag{1}
$$

where $A_{I_1 I_2}$ are the isotopic amplitudes, corresponding to the total isospin I_1 for the initial nucleons and the total isospin I_2 for the produced $\overline{D}N$ system.

One can see, from eq. (1), that seven different matrix elements for $N + N \rightarrow \Sigma_c + \overline{D} + N$ are characterized by three complex isotopic amplitudes: A_{11}, A_{10} and A_{01}. Therefore, in the general case, taking into account the possible interference between these amplitudes, one can find the following relation between the differential cross-sections:

$$
\frac{d\sigma}{d\omega}(pp \rightarrow \Sigma_c^{++} D^- p) + \frac{d\sigma}{drm\omega}(pp \rightarrow \Sigma_c^{++} \overline{D^0} n)
$$

$$
+ \frac{d\sigma}{d\omega}(np \rightarrow \Sigma_c^{++} D^- n) + \frac{d\sigma}{d\omega}(np \rightarrow \Sigma_c^{0} \overline{D^0} p) =
$$

$$
2 \left[\frac{d\sigma}{d\omega}(pp \rightarrow \Sigma_c^{+} \overline{D^0} p) + \frac{d\sigma}{d\omega}(np \rightarrow \Sigma_c^{+} D^- p) \right. \tag{2}
$$

$$
\left. + \frac{d\sigma}{d\omega}(np \rightarrow \Sigma_c^{+} \overline{D^0} n) \right],
$$

where $d\omega$ is the phase space element for the three-particle final state. Note that eq. (2) holds in any kinematical condition and for any reaction mechanism.

The situation essentially changes near the reaction threshold, where all final particles are produced in the s-state. The selection rules with respect to the Pauli principle, the P-invariance and the conservation of isospin and of the total angular momentum, allow to parametrize the spin structure of the matrix element for pp and np collisions as [10]

$$
\mathcal{M}(pp) = f_{10}(\tilde{\chi}_2 \, \sigma_y \, \sigma \cdot \hat{\mathbf{k}} \chi_1) \, (\chi_4^\dagger \sigma_y \, \tilde{\chi}_3^\dagger)
$$

$$
+ i f_{11}[\tilde{\chi}_2 \, \sigma_y(\sigma \times \hat{\mathbf{k}})_a \chi_1] \, (\chi_4^\dagger \sigma_a \sigma_y \tilde{\chi}_3^\dagger), \tag{3}
$$

$$
\mathcal{M}(np) = f_{01}(\tilde{\chi}_2 \, \sigma_y \chi_1) \, (\chi_4^\dagger \sigma \cdot \hat{\mathbf{k}} \sigma_y \tilde{\chi}_3^\dagger)
$$

$$
+ f_{10}(\tilde{\chi}_2 \sigma_y \sigma \cdot \hat{\mathbf{k}} \chi_1)(\chi_4^\dagger \sigma_y \tilde{\chi}_3^\dagger) \tag{4}
$$

$$
+ i f_{11}[\tilde{\chi}_2 \, \sigma_y(\sigma \times \hat{\mathbf{k}})_a \chi_1] \, (\chi_4^\dagger \sigma_a \sigma_y \tilde{\chi}_3^\dagger),
$$

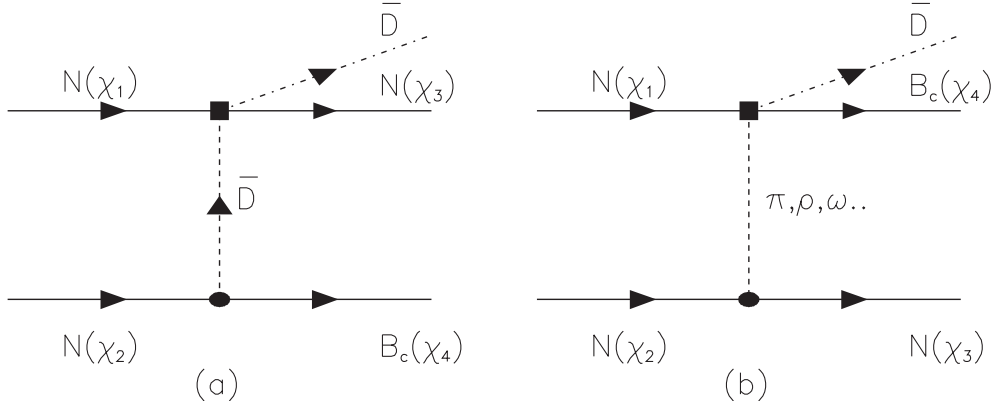**Fig. 1.** Feynman diagrams for the process $N + N \to \mathcal{B}_c + \overline{D} + N$, corresponding to two possible t exchanges: (a) $\overline{D}$-meson exchange, (b) light-mesons exchange.

where we are using the following notations: $\hat{\mathbf{k}}$ is the unit vector along the 3-momentum of the initial nucleon; χ_1 and χ_2 are the two-component spinors of the colliding nucleons, χ_3 and χ_4 are the two-component spinors of the scattered nucleon and of the produced charmed baryon; $\sigma_a = (\sigma_x, \sigma_y, \sigma_z)$ are the standard 2×2 Pauli matrices, f_{01} is the singlet-triplet amplitude for $\mathcal{B}_c = \Sigma_c$ or Λ_c threshold production, $f_{11}(f_{10})$ is the triplet-triplet (triplet-singlet) amplitude. One can see that the amplitude f_{01} is proportional to the isotopic amplitude A_{01}. Moreover, the isotopic amplitudes A_{11} and A_{10} have the same spin structure as in eq. (3). A direct consequence of the threshold spin structure, eqs. (3) and (4), is that the interferences $A_{11} \bigotimes A_{01}^*$ and $A_{10} \bigotimes A_{01}^*$, which can be present in the general case, vanish in the threshold region, after summing over the polarizations of the colliding nucleons. Therefore, we can derive the following relations, additional to (2), between the differential cross-sections, which hold for any model, at threshold:

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\omega}(np \to \Sigma_c^{++} D^- n) = \frac{\mathrm{d}\sigma}{\mathrm{d}\omega}(np \to \Sigma_c^0 \overline{D^0} p),$$

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\omega}(np \to \Sigma_c^+ D^- p) = \frac{\mathrm{d}\sigma}{\mathrm{d}\omega}(np \to \Sigma_c^+ \overline{D^0} n).$$

The normalization implicitly taken here is such that

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\omega} = |f_{01}|^2 + |f_{10}|^2 + 2|f_{11}|^2$$

for any process of np interaction, whereas the term $|f_{01}|^2$ is missing for the cross-section of any process of pp interaction.

3 Meson exchanges

Following the ideology of meson production in NN collisions, we consider different meson exchanges in t-channel, for different processes of associative charm particle production $N + N \to \mathcal{B}_c + \overline{D} + N$ (fig. 1). Taking into account the large difference in masses of the $\overline{D}$-meson and the light mesons, $\pi, \rho, \omega..$, one can expect mostly exchanges of

non-charmed mesons in t-channel. This is not the case for threshold production of heavy hadrons, where large values of momentum transfer squared are involved:

$$t_1 = (p_2 - p_4)^2 = -m(M - m) - \mu M \simeq -5.5 \text{ GeV}^2,$$

for the $\overline{D}$ exchange,

$$t_2 = (p_2 - p_3)^2 = -m(M - m) - \mu m \simeq -3.0 \text{ GeV}^2,$$

for the π exchange, where M, m, and μ are the masses of the charmed baryon, N and $\overline{D}$-meson, respectively. The ratio of the propagators, $(t_1 - \mu^2)/(t_2 - m_\pi^2) \simeq 3$, favors the π exchange, but the $\overline{D}$ contribution cannot be neglected. This kind of kinematical estimation does not work in case of associative strange-particle production, $N + N \to \Lambda(\Sigma) + K + N$, where, namely the K-meson exchange seems more probable than the π exchange [2]. This is supported by, at least, two arguments: the relative value of the total cross-section for Λ and Σ production near threshold, and the negative value of the D_{nn} coefficient in the process $\mathbf{p} + p \to \mathbf{\Lambda} + K^+ + p$ [11]. Therefore, as starting point for the present estimation, we will consider here only the $\overline{D}$ exchange.

Such model can be justified more rigorously using some quantitative arguments. The relative value of the D and π contributions to the matrix element of $N + N \to N + \mathcal{B}_c + \overline{D}$ processes, is determined not only by different π and D propagators, but also by the strength of the corresponding vertices. If the coupling constants $g_{N\mathcal{B}_c\overline{D}}$ and $g_{NN\pi}$ are comparable in absolute value, the amplitudes for $\overline{D} + N \to \overline{D} + N$ and $\pi + N \to \mathcal{B}_c + \overline{D}$ processes are essentially different. The first one is an elastic process, with non-vanishing threshold cross-section, whereas the second is an inelastic process, with a typical threshold decreasing. The high-energy behavior of the cross-sections is also different. Vector exchanges (ρ, ω mesons), which are mostly responsible for the scattering lengths of KN and $\overline{D}N$ elastic scattering, result in the same values: $a(KN) \simeq a(\overline{D}N)$. From ref. [12], one can find for the corresponding threshold amplitudes: $|f(K^+ p \to K^+ p)|^2 = |f(\overline{D}^0 p \to \overline{D}^0 p)|^2 = |a_1|^2 \simeq 910 \ \mu\text{b}$. At the same time,
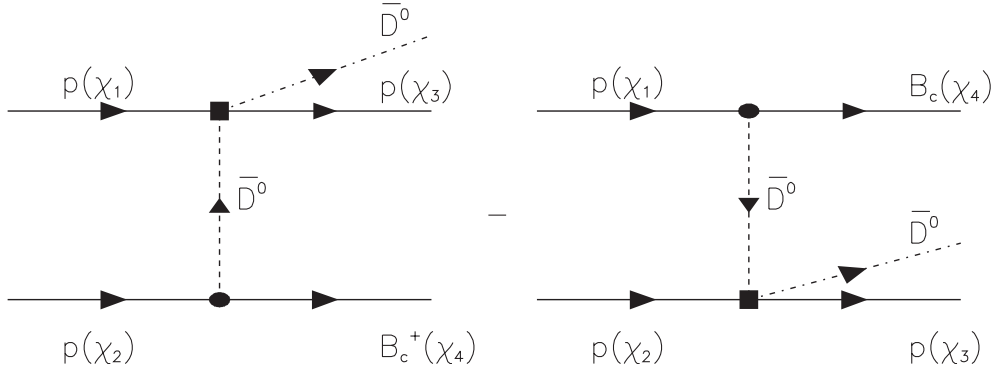
Fig. 2. $\overline{D^0}$ exchanges for the process $p + p \to \mathcal{B}_c + \overline{D^0} + p$.

the amplitude for $\pi^0 + p \to \Lambda + K^+$, (which is the $SU(4)$ equivalent of $\pi^0 + p \to \Lambda_c^+ + \overline{D^0}$), has a much smaller value: $|f(\pi^0 p \to \Lambda K^+)|^2 \simeq 30~\mu b$ [13].

A direct calculation of the cross-section for $\pi^- + p \to \Lambda_c^+ + D^-$[14] gives a value of 0.5 μb, at ($\sqrt{s} - \sqrt{s_0} = 10^{-3}$ GeV (s_0 is the threshold value of the total energy square s), whereas $\sigma(\overline{D}N \to \overline{D}N) \simeq 2000~\mu$b, *i.e.* three orders of magnitude larger.

These considerations are at the basis of the models which describe $p + p \to \Lambda + K^+ + p$ by the K-meson exchange, near threshold, which is confirmed by DISTO polarization transfer measurement [15]. Following $SU(4)$ symmetry, we can transpose these arguments to justify the D exchange in $\Lambda_c \overline{D}$ threshold production.

3.1 The reaction $p + p \to \Lambda_c^+(\Sigma_c^+) + \overline{D^0} + p$

Taking into account the identity of the colliding protons, the D exchange is described by two Feynman diagrams (fig. 2), with the following matrix element:

$$\mathcal{M}_D = \mathcal{M}_{1D} - \mathcal{M}_{2D},$$

$$\mathcal{M}_{1D} = -\frac{g_{p\mathcal{B}_c^+\overline{D^0}}}{t_1 - \mu^2}\frac{k\mathcal{N}}{E + m}\mathcal{A}(\overline{D^0}p \to \overline{D^0}p)$$
$$\times (\chi_3^\dagger \mathcal{I}\chi_1)(\chi_4^\dagger \sigma \cdot \hat{\mathbf{k}}\chi_2), \qquad (5)$$

$$\mathcal{M}_{2D} = -\mathcal{M}_{1D}~(\chi_1 \leftrightarrow \chi_2),$$

where $\mathcal{A}(\overline{D^0}p \to \overline{D^0}p)$ is the threshold amplitude for $\overline{D^0}p$ elastic scattering, $g_{p\mathcal{B}_c^+\overline{D^0}}$ is the pseudoscalar coupling constant for the vertex $p \to \mathcal{B}_c^+ + \overline{D^0}$, $\mathcal{N} = 2\sqrt{mM}(E+m)$ is the normalization factor, due to the presence of four baryons in the considered process, the factor $k/(E + m)$ results from the transformation of the Dirac spinor product $\overline{u}\gamma_5 u$ to its two-component equivalent, $\chi^\dagger \sigma \cdot \hat{\mathbf{k}}\chi$; E is the threshold energy of the initial proton in the reaction CMS, $\mathcal{I}$ is the unit 2×2 matrix.

We assume, in eq. (5), the pseudoscalar nature of the $N \to \mathcal{B}_c^+ + \overline{D^0}$ vertex, *i.e.* the product of the parities of all these particles must be $P(N\mathcal{B}_c^+\overline{D^0}) = -1$, following the quark model.

To transform the matrix element $\mathcal{M}_D$ into the "standard" form of eq. (4), we apply the Fierz transformation, in its two-component form:

$$(\chi_3^\dagger \mathcal{I}\chi_1)~(\chi_4^\dagger \sigma \cdot \hat{\mathbf{k}}\chi_2) = \frac{1}{2}\left\{-(\tilde{\chi}_2~\sigma_y~\chi_1)~(\chi_4^\dagger \sigma \cdot \hat{\mathbf{k}}\sigma_y\tilde{\chi}_3^\dagger)\right.$$
$$+ (\tilde{\chi}_2~\sigma_y\sigma \cdot \hat{\mathbf{k}}\chi_1)~(\chi_4^\dagger \sigma_y\tilde{\chi}_3^\dagger) \qquad (6)$$
$$\left.- i[\tilde{\chi}_2~\sigma_y(\sigma \times \hat{\mathbf{k}})_a\chi_1]~(\chi_4^\dagger \sigma_a\sigma_y\tilde{\chi}_3^\dagger)\right\}.$$

From eq. (6) one can see that $\mathcal{M}_{1D}$ contains not only the structures which are allowed by symmetry selection rules, but also a contribution which corresponds to a singlet-triplet transition (first term in eq. (6)). Such transition is forbidden by the generalized Pauli principle, following from the isotopic invariance of the strong interaction and should not appear in the total matrix element $\mathcal{M}_D$. This is an indication that only the contribution of both diagrams generates the correct answer. The expression for the total matrix element $\mathcal{M}_D$ is

$$\mathcal{M}_D(pp \to \mathcal{B}_c^+\overline{D^0}p) = \frac{g_{p\mathcal{B}_c^+\overline{D^0}}}{t_1 - \mu^2}\frac{k\mathcal{N}}{E + m}\mathcal{A}(\overline{D^0}p \to \overline{D^0}p)$$
$$\times \left\{-(\tilde{\chi}_2~\sigma_y~\sigma \cdot \hat{\mathbf{k}}\chi_1)~(\chi_4^\dagger \sigma_y\tilde{\chi}_3^\dagger)\right.$$
$$\left.+ i[\tilde{\chi}_2~\sigma_y(\sigma \times \hat{\mathbf{k}})_a\chi_1]~(\chi_4^\dagger \sigma_a\sigma_y\tilde{\chi}_3^\dagger)\right\}, \qquad (7)$$

with the necessary threshold spin structure and with definite relation between the threshold partial amplitudes:

$$f_{10}(pp \to \mathcal{B}_c^+\overline{D^0}p) = -f_{11}(pp \to \mathcal{B}_c^+\overline{D^0}p).$$

Comparing eqs. (6) and (7), one can conclude that the interference of the two diagrams (fig. 1) is very important, near threshold: a correct evaluation of all terms results in a 50% increase of the cross-section, compared with calculations based on non-interfering diagrams. A coherent consideration of the two contributions is especially important for the analysis of polarization phenomena in the near-threshold region.

3.2 The reaction $n + p \to \Lambda_c^+ + \overline{D^0} + n$

The $\overline{D}$ exchange, here, contains two Feynman diagrams — with exchange by neutral and charged D-mesons (fig. 3).
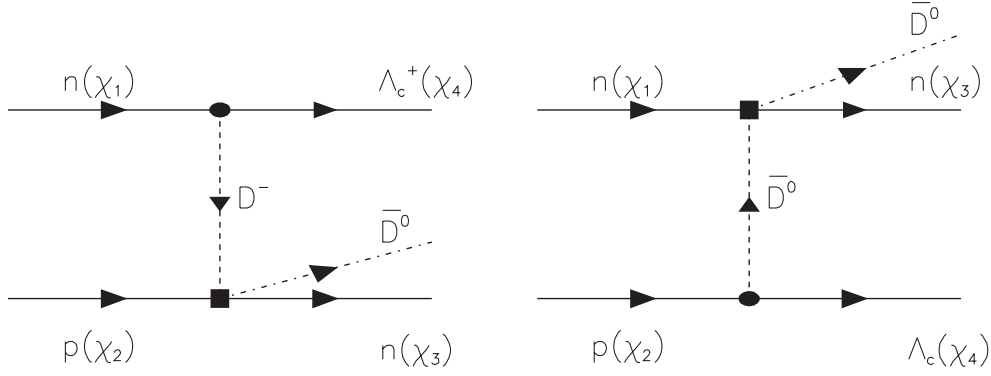
Fig. 3. D exchange for the process $n + p \rightarrow n + \Lambda_c + \overline{D^0}$.

The matrix element can be written as follows:

$$\mathcal{M}_D = \mathcal{M}_{\overline{D^0}} - \mathcal{M}_{D^-}.$$

Following the two-components Feynman rules for the considered diagrams, after the corresponding Fierz transformation, one can find

$$\mathcal{M}_{\overline{D^0}} = -\frac{g_{p\Lambda_c^+\overline{D^0}}}{t_1 - \mu^2}\frac{k\mathcal{N}}{E + m}\mathcal{A}(\overline{D^0}n \rightarrow \overline{D^0}n)$$
$$\times(\chi_4^\dagger\sigma \cdot \hat{\mathbf{k}}\chi_2)(\chi_3^\dagger\mathcal{I}\chi_1) =$$
$$\frac{1}{2}\frac{g_{p\Lambda_c^+\overline{D^0}}}{t_1 - \mu^2}\frac{k\mathcal{N}}{E + m}\mathcal{A}(\overline{D^0}n \rightarrow \overline{D^0}n)$$
$$\times\left\{(\tilde{\chi}_2\ \sigma_y\ \chi_1)\ (\chi_4^\dagger\sigma \cdot \hat{\mathbf{k}}\sigma_y\tilde{\chi}_3^\dagger)\right.$$
$$-(\tilde{\chi}_2\ \sigma_y\sigma \cdot \hat{\mathbf{k}}\chi_1)\ (\chi_4^\dagger\sigma_y\tilde{\chi}_3^\dagger)$$
$$\left. + i[\tilde{\chi}_2\ \sigma_y(\sigma \times \hat{\mathbf{k}})_a\chi_1]\ (\chi_4^\dagger\sigma_a\sigma_y\tilde{\chi}_3^\dagger)\right\}. \quad (8)$$

Taking into account that $\mathcal{M}_{D^-}$ can be derived from $\mathcal{M}_{\overline{D^0}}$, after the following substitutions:

$$\chi_1 \leftrightarrow \chi_2,$$
$$\hat{\mathbf{k}} \leftrightarrow -\hat{\mathbf{k}},$$
$$\mathcal{A}(\overline{D^0}n \rightarrow \overline{D^0}n) \leftrightarrow \mathcal{A}(D^-p \rightarrow \overline{D^0}n),$$

one can obtain the following formula, for the total matrix element:

$$\mathcal{M}_D(np \rightarrow p\Lambda_c^+\overline{D^0}) = \frac{1}{2}\frac{g_{p\Lambda_c^+\overline{D^0}}}{t_1 - \mu^2}\frac{k\mathcal{N}}{E + m}$$
$$\times\left\{(\tilde{\chi}_2\ \sigma_y\ \chi_1)\ (\chi_4^\dagger\sigma \cdot \hat{\mathbf{k}}\sigma_y\tilde{\chi}_3^\dagger)\right.$$
$$\times\left[\mathcal{A}(\overline{D^0}n \rightarrow \overline{D^0}n) - \mathcal{A}(D^-p \rightarrow \overline{D^0}n)\right]$$
$$+\left[-(\tilde{\chi}_2\ \sigma_y\sigma \cdot \hat{\mathbf{k}}\ \chi_1)\ (\chi_4^\dagger\sigma_y\tilde{\chi}_3^\dagger)\right.$$
$$\left. + i\left(\tilde{\chi}_2\ \sigma_y(\sigma \times \hat{\mathbf{k}})_a\chi_1\right)(\chi_4^\dagger\sigma_a\sigma_y\tilde{\chi}_3^\dagger)\right]$$
$$\left.\times\left[\mathcal{A}(\overline{D^0}n \rightarrow \overline{D^0}n) + \mathcal{A}(D^-p \rightarrow \overline{D^0}n)\right]\right\}. \quad (9)$$

Using the isotopic invariance for the processes of $\overline{D}N$ scattering,

$$\mathcal{A}(\overline{D^0}n \rightarrow \overline{D^0}n) - \mathcal{A}(D^-p \rightarrow \overline{D^0}n) = a_0,$$
$$\mathcal{A}(\overline{D^0}n \rightarrow \overline{D^0}n) + \mathcal{A}(D^-p \rightarrow \overline{D^0}n) = a_1,$$

where a_0 and a_1 are the $\overline{D}N$-scattering lengths, corresponding to the possible values of the total isospin for $\overline{D}N$ scattering, one can see that the f_{01}-amplitude for $n + p \rightarrow \Lambda_c^+ + \overline{D^0} + n$ depends only on a_0, and both triplet amplitudes, f_{10} and f_{11}, depend on a_1 —in agreement with the correlation between the spin and the isospin structures discussed above. Note that, despite the large mass of the D-meson, the scattering lengths a_0 and a_1 can be considered real quantities, as in case of KN scattering, for the same physical reason: following the unitarity condition there is no intermediate state for the $\overline{D}N$ channel at zero $\overline{D}$-meson energy, due to the absence of baryons with negative charm.

The expressions (7) and (9) allow to write the two isotopic amplitudes A_1 and A_0, for $\Lambda_c^+\overline{D^0}$ production in NN collisions:

$$A_1 = a_1\frac{g_{p\Lambda_c^+\overline{D^0}}}{t_1 - \mu^2}\frac{k\mathcal{N}}{E + m}\left\{-(\tilde{\chi}_2\ \sigma_y\sigma \cdot \hat{\mathbf{k}}\ \chi_1)\ (\chi_4^\dagger\sigma_y\tilde{\chi}_3^\dagger)\right.$$
$$\left. + i\left[\tilde{\chi}_2\ \sigma_y(\sigma \times \hat{\mathbf{k}})_a\chi_1\right](\chi_4^\dagger\sigma_a\sigma_y\tilde{\chi}_3^\dagger)\right\},$$
$$A_0 = a_0\frac{g_{p\Lambda_c^+\overline{D^0}}}{t_1 - \mu^2}\frac{k\mathcal{N}}{E + m}(\tilde{\chi}_2\ \sigma_y\chi_1)(\chi_4^\dagger\sigma \times \hat{\mathbf{k}}\sigma_y\tilde{\chi}_3^\dagger).$$

These expressions allow us to find the ratio $\mathcal{R}_D$ of the cross-sections of $\Lambda_c\overline{D}$ production in np and pp collisions:

$$\mathcal{R}_D = \frac{\sigma(np \rightarrow \Lambda_c^+\overline{D^0}n)}{\sigma(pp \rightarrow \Lambda_c^+\overline{D^0}p)} = \frac{\sigma(np \rightarrow \Lambda_c^+D^-p)}{\sigma(pp \rightarrow \Lambda_c^+\overline{D^0}p)} =$$
$$\frac{1}{4}\left(1 + \frac{a_0^2}{3a_1^2}\right) > \frac{1}{4}, \quad (10)$$

where the index D underlines that this prediction is valid only in the framework of the D exchange.

This result can be modified by the presence of the Λ_cN interaction in the S-state. Taking into account the singlet

$a_s(\Lambda_c N)$ and triplet $a_t(\Lambda_c N)$ scattering lengths, a simple estimation of this effect gives

$$\mathcal{R}_D^{\text{FSI}} = \mathcal{R}_D + \frac{1}{6}(1 - x_{ts})\frac{a_0^2}{a_1^2}, \qquad (11)$$

where $x_{ts} = a_t^2(\Lambda_c N)/a_s^2(\Lambda_c N)$, i.e. FSI will decrease (increase) the ratio $\mathcal{R}_D$ when the triplet $\Lambda_c N$ scattering is stronger (weaker) than the singlet $\Lambda_c N$ scattering. In any case, $\mathcal{R}$ is sensitive to FSI interaction. In the initial state, taking into account that the reaction threshold is quite large, the pp and np interactions can be assumed similar. Therefore, one expects that ISI, for the calculation of $\mathcal{R}_D$, are less important than in the case of strange-particle production. Note that $|a_0/a_1| \simeq 0.1$ for elastic KN scattering, and it is likely that this estimation is correct also for elastic $\overline{D}N$ scattering. In this context the ratio (11) will be insensitive to $\Lambda_c N$ FSI interaction too.

4 Estimation of the cross-section

We give here an estimation of the cross-section for $N + N \to \Lambda_c + \overline{D} + N$ in the threshold region, where the energy dependence of the total cross-section can be written as [2]

$$\sigma_{pp \to \Lambda_c DN} = \overline{|\mathcal{M}_D|^2}V_{ps}/\mathcal{F}, \qquad (12)$$

with

$$V_{ps} = \frac{\pi^3}{2}\frac{(mM\mu)^{1/2}}{(m + M + \mu)^{3/2}}Q^2,$$

$$Q = \sqrt{s} - (m + M + \mu), \ \mathcal{F} = 4\sqrt{s}k(2\pi)^5, \qquad (13)$$

where we assumed, for simplicity, that $\overline{|\mathcal{M}|_D^2}$ is constant, in the region of applicability of Ss production of final particles.

Let us consider the ratio of the threshold matrix elements for the processes $p + p \to \Lambda_c^+ + \overline{D^0} + p$ (D exchange) and $p + p \to \Lambda + K^+ + p$ (K exchange). One can find

$$\mathcal{R}_{DK} = \frac{\overline{|\mathcal{M}_D|^2}}{\overline{|\mathcal{M}_K|^2}} = \left(\frac{g_{p\Lambda_c^+\overline{D^0}}}{g_{p\Lambda K^+}}\right)^2\left[\frac{a_1(DN)}{a_1(KN)}\right]^2$$

$$\times \left(\frac{t_K - m_K^2}{t_D - m_D^2}\right)^2\frac{M(\Lambda_c)}{M(\Lambda)}\frac{(E_D^2 - m^2)}{(E_K^2 - m^2)}, \qquad (14)$$

where $E_{D,K} = \frac{1}{2}(M + m + \mu)$, $M(\mu)$ is the mass of $\Lambda_c(D)$ for D production and $M(\mu)$ is the mass of $\Lambda(K)$ for K-meson production, t_K and t_D are the corresponding momentum transfer squared in K- or D-meson threshold production. Taking the numerical values

$$M(\Lambda_c)/M(\Lambda) \simeq 2.05, \ (E_D^2 - m^2)/(E_K^2 - m^2) = 7.53,$$

$$(t_K - m_K^2)^2/(t_D - m_D^2)^2 \simeq 1.1 \cdot 10^{-2}, \qquad (15)$$

one can find

$$\mathcal{R}_{DK} = 0.11\left(\frac{g_{p\Lambda_c^+\overline{D^0}}}{g_{p\Lambda K^+}}\right)^2\left[\frac{a_1(DN)}{a_1(KN)}\right]^2. \qquad (16)$$

Therefore, the ratio of total cross-sections, at a fixed value of Q, for $\Lambda_c^+\overline{D^0}$ and ΛK^+ production in pp collisions can be written as

$$\frac{\sigma_t(pp \to \Lambda_c^+\overline{D^0}p)}{\sigma_t(pp \to \Lambda K^+p)} = \mathcal{R}_{DK}\left[\frac{M(\Lambda_c)\mu(\overline{D^0})}{M(\Lambda)\mu(K)}\right]^{1/2}$$

$$\times \left[\frac{m + M(\Lambda) + \mu(K)}{m + M(\Lambda_c) + \mu(\overline{D^0})}\right]^{5/2}\frac{k(\Lambda)}{k(\Lambda_c)} =$$

$$0.27\mathcal{R}_{DK} = 2.98 \cdot 10^{-2}\left(\frac{g_{p\Lambda_c^+\overline{D^0}}}{g_{p\Lambda K^+}}\right)^2\left[\frac{a_1(DN)}{a_1(KN)}\right]^2. \quad (17)$$

We neglected, in this estimation, an important ingredient, the phenomenological hadron form factors, which are traditionally included in these calculations. Such form factors will suppress the charmed-particle production near threshold, according to

$$\mathcal{D}_F = \left(1 - \frac{t_K - m_K^2}{\lambda^2}\right)^2 \bigg/ \left(1 - \frac{t_D - m_D^2}{\lambda^2}\right)^2, \quad (18)$$

where, for simplicity, each hadronic vertex is described by the same monopole form factor, with the same cut-off parameter λ. Taking, as an example, $\lambda = 1.5$ GeV, one can find a suppression factor $\mathcal{D}_F \simeq 0.08$.

In the framework of the SU_4 symmetry, the coupling constants $g_{p\Lambda_c^+\overline{D^0}}$ and $g_{p\Lambda K^+}$ as well as the scattering lengths $a_1(DN)$ and $a_1(KN)$ are equal, so finally,

$$\frac{\sigma_t(pp \to \Lambda_c^+\overline{D^0}p)}{\sigma_t(pp \to \Lambda K^+p)} \simeq 3 \cdot 10^{-2}. \qquad (19)$$

Dipole hadronic form factors will result in a cross-section one order of magnitude smaller, for $\Lambda_c^+\overline{D^0}$ production. On the basis of eq. (19), taking the experimental data on the total cross-section for $p + p \to \Lambda + K^+ + p$ [2], one can predict the following energy dependence for the cross-section of open charm production:

$$\sigma(pp \to \Lambda_c^+\overline{D^0}p) = 0.1\text{--}1(Q/0.1 \text{ GeV})^2 \text{ nb}, \qquad (20)$$

depending on the choice of form factors. As we mentioned above, no experimental data exist in the threshold region. The lowest energy where the open charm cross-section has been measured, $E_p = 70$ GeV, $\sqrt{s} = 11.46$ GeV, corresponds to $Q \simeq 6$ GeV. From eq. (20) we find, for this energy: $\sigma(pp \to \Lambda_c^+\overline{D^0}p) \simeq 0.7$ μb, to be compared to the experimental value $\sigma(pp \to$ charm$) = 1.6^{+1.1}_{-0.7}$ (stat) ± 0.3 (syst) μb [4]. However, this apparent agreement, resulting form a very large extrapolation in Q, can be accidental, because eqs. (12) and (13) rigorously hold only for $Q < \mu$.

5 Conclusions

In this paper we analyzed the simplest processes of open charm production in the near-threshold region of NN collisions, $N + N \to \Lambda_c(\Sigma_c) + \overline{D} + N$, in a model-independent

way, using the symmetry properties of the strong interaction, in analogy with strange-particle production [10]. We established the spin and isospin structure of the corresponding threshold matrix element in exact form and showed a strong correlation of spin and isospin variables.

The D exchange mechanism for $N + N \rightarrow \Lambda_c(\Sigma_c) + \overline{D} + N$ can be considered a natural generalization of the K exchange for strange-particle production, which appears as the main mechanism in the near-threshold region.

We have described the open charm production in terms of the following ingredients:

- the coupling constants $g_{N\Lambda_c\overline{D}}$ and $g_{N\Sigma_c\overline{D}}$ and the threshold amplitude for the elementary $\overline{D}N$ scattering (both these quantities can be related to the corresponding values for strange particles, namely the $g_{N\Lambda K}$ coupling constant and the KN-scattering length, through SU_4 symmetry);
- the parameters (scattering lengths and effective radii) of final strong $N\Lambda_c$ interaction described in analogy with low-energy ΛN interaction;
- the hadronic phenomenological form factors.

This simple and transparent model allows us to predict the relative value of the cross-sections for $\Lambda_c\overline{D}$ and ΛK production, at the same value of the energy excess Q. We find a difference of three orders of magnitude in the values of the corresponding cross-sections, on the basis of simple kinematical arguments, such as the difference of the reaction thresholds and of the meson propagators (due to the difference of the threshold momentum transfer). Phenomenological hadron form factors also play a large role in the threshold region of charm production.

References

1. H. Feldmeier, J. Knoll, W. Norenberg, J. Wambach (Editors), *Proceedings of Hirschegg '01: Structure of Hadrons: 29th International Workshop on Gross Properties of Nuclei and Nuclear Excitations, Hirschegg, Austria, 14-20 January 2001* (GSI, Darmstadt, 2001) p. 370.
2. P. Moskal, M. Wolke, A. Khoukaz, W. Oelert, Prog. Part. Nucl. Phys. **49**, 1 (2002).
3. E. Botta, T. Bressani, A. Feliciello (Editors), *Proceedings of HYP2000: 7th International Conference on Hypernuclear and Strange Particle Physics, Torino, Italy, 23-27 October 2000* (North-Holland, Amsterdam, 2001), in Nucl. Phys. A **691** (2001), Vols. **1**, **2**.
4. SVD Collaboration (N.S. Amaglobeli *et al.*), Phys. Atom. Nucl. **64**, 891 (2001).
5. K.G. Boreskov, A.B. Kaidalov, Sov. J. Nucl. Phys. **37**, 100 (1983).
6. NA50 Collaboration (M.C. Abreu *et al.*), Eur. Phys. J. C **14**, 443 (2000).
7. W. Cassing, L.A. Kondratyuk, G.I. Lykasov, M.V. Rzjanin, Phys. Lett. B **513**, 1 (2001).
8. NA50 Collaboration (P. Bordalo *et al.*), Nucl. Phys. A **661**, 538 (1999).
9. W. Cassing, E.L. Bratkovskaya, A. Sibirtsev, Nucl. Phys. A **691**, 753 (2001).
10. M.P. Rekalo, J. Arvieux, E. Tomasi-Gustafsson, Phys. Rev. C **56**, 2238 (1997).
11. J.M. Laget, Nucl. Phys. A **691**, 11c (2001).
12. M. Hoffmann, J.W. Durso, K. Holinde, B.C. Pearce, J. Speth, Nucl. Phys. A **593**, 341 (1995).
13. A. Gasparian, J. Haidenbauer, C. Hanhart, L. Kondratyuk, J. Speth, Phys. Lett. B **480**, 273 (2000).
14. A.M. Gasparyan, V.Y. Grishina, L.A. Kondratyuk, W. Cassing, J. Speth, arXiv:nucl-th/0210018.
15. DISTO Collaboration (F. Balestra *et al.*), Phys. Rev. Lett. **83**, 1534 (1999).