Spin Observables in Transition-Distribution-Amplitude Studies

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Abstract. Exclusive hadronic reactions with a massive lepton pair ($\ell^+\ell^-$) in the final state will be measured with PANDA at GSI-FAIR and with Compass at CERN, both in $p + \bar{p} \rightarrow \ell^+\ell^- + \pi$ and $\pi + N \rightarrow N' + \ell^+\ell^-$. Similarly, electroproduction of a meson in the backward region will be studied at JLAB. We discuss here how the spin structure of the amplitude for such processes will enable us to disentangle various mechanisms. For instance, target-transverse-spin asymmetries are specific of a partonic description, where the amplitude is factorised in terms of baryon to meson or meson to baryon Transition Distribution Amplitudes (TDAs) as opposed to what is expected from baryon-exchange contributions.

1. Introduction

In [1, 2], we have introduced the concept of Transition Distribution Amplitudes (TDAs) containing new information on the hadron structure. These non-perturbative objects appear in the study of exclusive reactions in a new scaling regime, i.e. involving a large –timelike or spacelike– Q^2 photon and a baryonic exchange in the t-channel. This extends the concept of Generalised Parton Distributions (GPDs), as already advocated in the pioneering work of [3].

In particular, we have discussed [4] the backward electroproduction of a pion,

$$\gamma^\star(q)N(p_1) \rightarrow N'(p_2)\pi(p_\pi), \tag{1}$$

on a proton (or neutron) target, whose amplitude can be factorised in a hard part convoluted with TDAs (see Fig. 1 (a)). Along the same lines, the reaction,

$$N(p_1)\bar{N}(p_2) \rightarrow \gamma^\star(q)\pi(p_\pi), \tag{2}$$

in the near forward region was also studied in terms of TDAs [5, 6, 7] (see Fig. 1 (b)). The TDAs involved in the description of Deeply-Virtual Compton Scattering (DVCS) in the backward kinematics $\gamma^\star(q)N(p_1) \rightarrow N'(p_2)\gamma(p_\gamma)$ and the reaction $N(p_1)\bar{N}(p_2) \rightarrow \gamma^\star(q)\gamma(p_\gamma)$ in the near forward region were given in [8].

Recently, two studies of the proton-to-pion TDAs were carried out, one [9] in the meson-cloud model, another [10] guided by the concept of the spectral representation [11, 12]. Yet, more work is needed before being able to proceed to *quantitative* comparisons between different TDA models and

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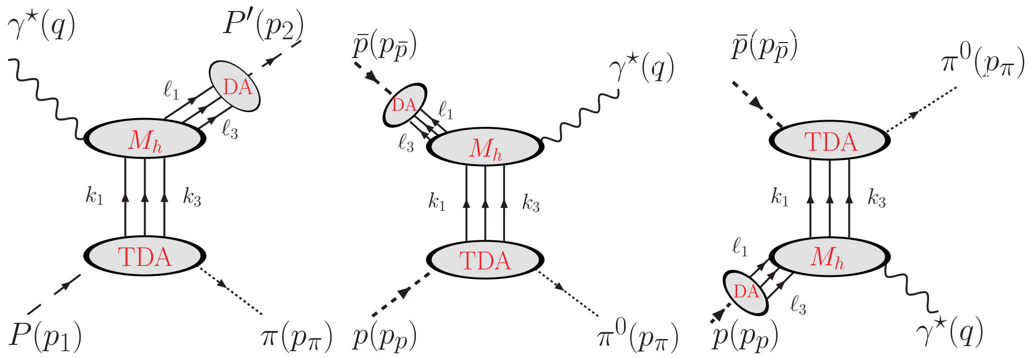


(a) π^0 electroproduction at small u (b) $p\bar{p} \to \gamma^\star \pi^0$ at small $t = (p_p - p_{\pi^0})^2$ (c) $p\bar{p} \to \gamma^\star \pi^0$ at small $u = (p_{\bar{p}} - p_{\pi^0})^2$

Figure 1. Illustration of the factorisation for three exclusive reactions involving the TDAs.

between theory and experiments. For the time being, model independent analyses sound more expedient. These can be divided in three categories: Q^2 dependence and scaling analyses, dominance of specific polarisations of the off-shell photon and target (or projectile) transverse-spin asymmetries. In this note, we shall review these.

2. Q^2-dependence, scaling and dominant photon polarisation: Example of the TDA studies in $\bar{p}p$ annihilation

Let us analyse here the case of $\bar{p}p \to \ell^+\ell^-\pi$. This will enable us to discuss the scaling and Q^2-dependence in the factorised picture as well as the dominance of specific photon polarisations. The momenta of the subprocess $\bar{p}p \to \gamma^\star\pi$ are defined as shown in Fig. 1 (b). We shall first limit our discussion to the region where $t = (p_p - p_\pi)^2$ is much smaller than the invariant mass of the $\ell^+\ell^-$ pair, Q^2. The region where $u = (p_{\bar{p}} - p_\pi)^2 \ll Q^2$ (see Fig. 1 (c)) is briefly discussed at the end of the section.

2.1. Kinematics

The z-axis is chosen along the colliding proton and anti-proton and the $x - z$ plane is identified with the collision or hadronic plane. We define the light-cone vectors p and n such that $2\ p.n = 1$, as well as $P = (p_p + p_\pi)/2$, $\Delta = p_\pi - p_p$ and its transverse component Δ_T ($\Delta_T^2 < 0$). ξ is defined as $\xi = -\frac{\Delta.n}{2P.n}$. We express the particle momenta through a Sudakov decomposition, which, for $\Delta_T = 0$, $M \ll W$ and $m_\pi = 0$, is

$$p_p = (1+\xi)p, \quad p_{\bar{p}} = \frac{W^2}{1+\xi}n, \quad p_\pi = (1-\xi)p, \quad t = \frac{2\xi M^2}{1+\xi}, \quad \xi = \frac{Q^2}{2W^2 - Q^2}. \tag{3}$$

In the fixed-target mode, the maximal reachable value for $W^2 = 2M^2 + 2ME_{\bar{p}}$ at GSI will be $\simeq 30$ GeV^2 (for $E_{\bar{p}} = 15$ GeV). The highest invariant mass of the photon could be $Q^2_{max} \simeq 30$ GeV^2. We refer to [13] for a complete discussion of the kinematically allowed domain. In terms of our notations, in the proton rest frame, we have $p = \frac{M}{2(1+\xi)}(1,0,0,-1)$ and $n = \frac{1+\xi}{2M}(1,0,0,1)$. Thus $\xi \in [0.5, 1]$ corresponds to $|p_\pi^z| < M/6 \simeq 155$ MeV in the laboratory frame at $\Delta_T = 0$.

2.2. The properties of the amplitude

At $\Delta_T = 0$, the leading-twist TDAs for the $p \to \pi^0$ transition, $V_i^{p\pi^0}(x_i,\xi,\Delta^2)$, $A_i^{p\pi^0}(x_i,\xi,\Delta^2)$ and $T_i^{p\pi^0}(x_i,\xi,\Delta^2)$ are defined as:

$$\mathcal{F}\Big(\langle\pi^0(p_\pi)|\,\epsilon^{ijk}u^i_\alpha(z_1n)u^j_\beta(z_2n)d^k_\gamma(z_3n)\,|P(p_p,s_p)\rangle\Big) = \frac{i}{4}\frac{f_N}{f_\pi}\Big[V_1^{p\pi^0}(\not{p}C)_{\alpha\beta}(u^+(p_p,s_p))_\gamma$$
$$+A_1^{p\pi^0}(\not{p}\gamma^5C)_{\alpha\beta}(\gamma^5u^+(p_p,s_p))_\gamma + T_1^{p\pi^0}(\sigma_{p\mu}C)_{\alpha\beta}(\gamma^\mu u^+(p_p,s_p))_\gamma\Big]\,,$$

where $\sigma^{\mu\nu} = 1/2[\gamma^\mu, \gamma^\nu]$, C is the charge conjugation matrix, $f_\pi = 131$ MeV is the pion decay constant and $f_N \sim 5.2 \cdot 10^{-3}$ GeV^2. u^+ is the large component of the nucleon spinor such that one has $u(p_p, s_p) = (\not{n}\not{p} + \not{p}\not{n})u(p_p, s_p) = u^-(p_p, s_p) + u^+(p_p, s_p)$ with $u^+(p_p, s_p) \sim \sqrt{p_p^+}$ and $u^-(p_p, s_p) \sim \sqrt{1/p_p^+}$.

At the leading order in α_s and at $\Delta_T = 0$, the amplitude $\mathcal{M}_\lambda^{s_p s_{\bar{p}}}$ for $\bar{p}(p_{\bar{p}}, s_{\bar{p}})p(p_p, s_p) \to \gamma^\star(q, \lambda)\pi^0(p_\pi)$ reads

$$\mathcal{M}_\lambda^{s_p s_{\bar{p}}} = -i\frac{(4\pi\alpha_s)^2\sqrt{4\pi\alpha_{em}}f_N^2}{54 f_\pi Q^4} \underbrace{\bar{v}^+(p_{\bar{p}}, s_{\bar{p}})\not{\epsilon}^\star(\lambda)\gamma^5 u^+(p_p, s_p)}_{\mathcal{S}_\lambda^{s_p s_{\bar{p}}}} \underbrace{\int\limits_{-1+\xi}^{1+\xi}[dx]\int\limits_0^1[dy]\left(2\sum_{\alpha=1}^{7}R_\alpha + \sum_{\alpha=8}^{14}R_\alpha\right)}_{\mathcal{I}}, \quad (4)$$

where $[dx] = dx_1dx_2dx_3\delta(2\xi - \sum_k x_k)$ and $[dy] = dy_1dy_2dy_3\delta(1-\sum_k y_k)$; the coefficients R_α ($\alpha = 1, ..., 14$) exactly correspond to T_α in [4] after the replacement $-i\epsilon \to i\epsilon$ due to the presence of the $\gamma^\star$ in the final instead of initial state.

Denoting the linear polarisations of the virtual photon by the indices L, x, y, one defines [14, 15] $\sigma_T \propto 1/2[\mathcal{M}_x(\mathcal{M}_x)^* + \mathcal{M}_y(\mathcal{M}_y)^*]$, $\sigma_L \propto \mathcal{M}_L(\mathcal{M}_L)^*$, $\sigma_{TL} \propto \mathcal{M}_x(\mathcal{M}_L)^* + \mathcal{M}_L(\mathcal{M}_x)^*$ and $\sigma_{TT} \propto 1/2[\mathcal{M}_x(\mathcal{M}_x)^* - \mathcal{M}_y(\mathcal{M}_y)^*]$. The corresponding definitions apply for the squared of $\mathcal{S}_\lambda^{s_p s_{\bar{p}}}$ summed over the proton spins, $\mathcal{S}^2_{T,L,LT,TT}$.

At the leading twist, only $\mathcal{S}_T^2 = (2(1+\xi)Q^2)/\xi$ survives. This means that the far off-shell photon produced in association with the pion is dominantly transversely polarised. This dominance increases with Q^2. As a direct consequence, the angular dependence of the lepton pair follows the distribution

$$1 + \cos^2\theta_\ell, \quad (5)$$

where θ_ℓ is the angle between the ℓ^+ momentum in the dilepton rest frame and the z axis. Along the same line of argument, there should not be any azimuthal dependence, ϕ_ℓ, in the dilepton distribution. We should stress here that, at $|\Delta_T|^2 = 0$, these results coincide with the ones for the proton form factor studies in the timelike region. However, for non-vanishing transverse momenta, the spin-quantisation axis may be rotated inducing different definitions of the aforementioned angles.

Let us now discuss the scaling of the amplitude squared. While the contribution from $\mathcal{I}$ may depend on ξ (and Δ_T^2), it does not depend on Q^2 (except for a logarithmic dependence of the TDAs if one takes into account their QCD evolution). Given that $\mathcal{S}_T^2 = (2(1+\xi)Q^2)/\xi$, we deduce from Eq. (4) that the amplitude squared scales like Q^{-6} at fixed ξ.

2.3. The $\bar{p} \to \pi^0$ transition region

So far, we have discussed the case where the momentum-transfer squared between the proton and the pion is small compared to Q^2, *i.e.* when the pion is moving slowly in the proton-rest frame. However, nothing prevents us for applying the same arguments for the factorisation of the amplitude in the case where the pion is slow in the anti-proton rest frame. One expects the same kind of mechanism to take place. The amplitude factorises into a hard part and anti-proton-to-pion TDAs. In practice, at GSI-FAIR, one would observe a very energetic (and thus near forward) pion and a very slow lepton pair in the laboratory frame. The same scaling properties of the amplitude and the same expectations for the dilepton polarisation would obviously hold.

3. Single Transverse Spin Asymmetry: Example of the TDA studies in backward electroproduction of a pion

In order to study a possible Single Spin Asymmetry (SSA) in backward electroproduction of a pion on a transversely polarised target, we shall study the quantity $\sigma^{s_1} - \sigma^{-s_1}$ with the definition

$$\sigma^{s_1} = \sum_\lambda \sum_{s_2} (\mathcal{M}_\lambda^{s_1 s_2})(\mathcal{M}_\lambda^{s_1 s_2})^*. \quad (6)$$

As we shall see the SSA can only arise for $\Delta_T \neq 0$ and will be a function of azimuthal angles of the process. We thus need to define precisely the kinematics and say some words about $\mathcal{M}_\lambda^{s_1 s_2}$ for $\Delta_T \neq 0$.

3.1. Kinematics

The momenta of the process $\gamma^\star P \to P'\pi$ are defined as in Fig. 1 (a) and Fig. 2. The z-axis is chosen along the initial-nucleon and the virtual-photon momenta. The $x - z$ plane is identified with the collision or hadronic plane (Fig. 2). Here, $P = \frac{1}{2}(p_1 + p_\pi)$, $\Delta = p_\pi - p_1$ and its transverse component is Δ_T ($\Delta_T.\Delta_T = \Delta_T^2 < 0$).

We can then express the momenta of the particles through their Sudakov decomposition and, keeping the first-order corrections in the masses and Δ_T^2, we have:

$$\begin{aligned} p_1 &= (1+\xi)p + \frac{M^2}{1+\xi}n, q \simeq -2\xi\Big(1+\tfrac{(\Delta_T^2-M^2)}{Q^2}\Big)p + \tfrac{Q^2}{2\xi}\Big(1+\tfrac{(\Delta_T^2-M^2)}{Q^2}\Big)^{-1} n, \\ p_\pi &= (1-\xi)p + \frac{m_\pi^2-\Delta_T^2}{1-\xi}n + \Delta_T, \Delta = -2\xi p + \Big[\frac{m_\pi^2-\Delta_T^2}{1-\xi} - \frac{M^2}{1+\xi}\Big]n + \Delta_T \\ p_2 &\simeq -2\xi\frac{(\Delta_T^2-M^2)}{Q^2}p + \Big[\frac{Q^2}{2\xi}\Big(1+\tfrac{(\Delta_T^2-M^2)}{Q^2}\Big)^{-1} - \tfrac{m_\pi^2-\Delta_T^2}{1-\xi} + \tfrac{M^2}{1+\xi}\Big]n - \Delta_T, \end{aligned} \tag{7}$$

with $\xi \simeq \frac{Q^2}{Q^2+2(W^2+\Delta_T^2-M^2)}$.

For $\varepsilon_x = (0,1,0,0)$ and $\varepsilon_y = (0,0,1,0)$ with the axis definitions of Fig. 2, one may further specify that

$$\Delta_T = |\Delta_T|(\cos\phi\,\varepsilon_x + \sin\phi\,\varepsilon_y) \text{ and } s_{T,1} = s_1 = \cos\phi_S\,\varepsilon_x + \sin\phi_S\,\varepsilon_y, \tag{8}$$

for the transverse spin of the target ($s_1.p_1 = s_1.p = s_1.n = 0$).

θ_π^* is defined as the polar angle between the virtual photon and the pion in the $P'\pi^0$ center-of-mass frame (see Fig. 2). ϕ is the azimuthal angle between the electron plane and the plane of the process $\gamma^\star P \to P'\pi^0$ (hadronic plane) ($\phi = 0$ when the pion is emitted in the half plane containing the outgoing electron).

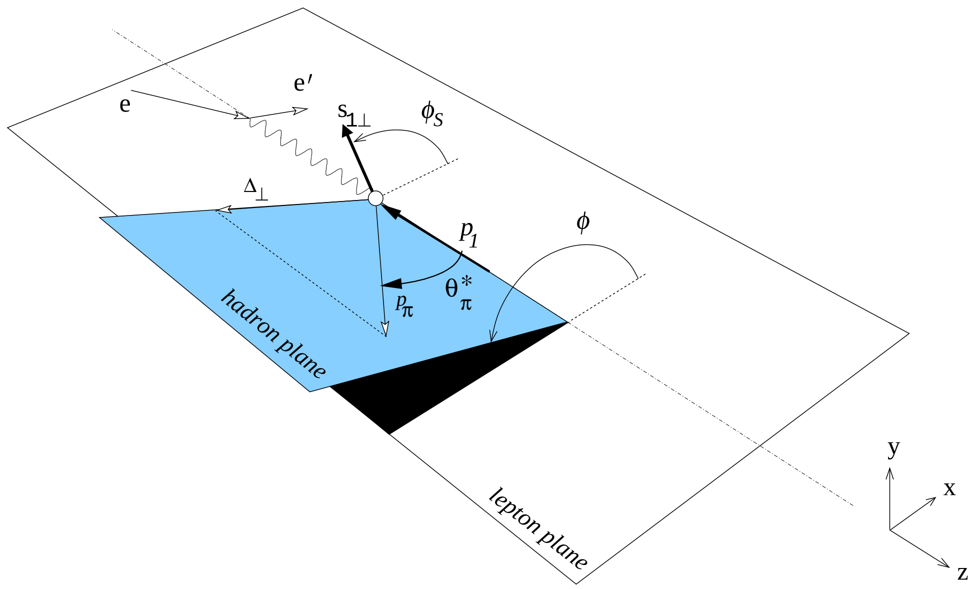


Figure 2. Kinematics of electroproduction of a pion and definition of the angles ϕ and ϕ_S

3.2. Hard-amplitude at $\Delta_T \neq 0$

At leading order in α_s, the amplitude $\mathcal{M}_\lambda^{s_1 s_2}$ for $\gamma^\star(q,\lambda)P(p_1,s_1) \to P'(p_2,s_2)\pi^0(p_\pi)$ reads

$$\mathcal{M}_\lambda^{s_1 s_2} = \underbrace{-i\frac{(4\pi\alpha_s)^2\sqrt{4\pi\alpha_{em}}f_N^2}{54 f_\pi}}_{C}\frac{1}{Q^4}\Big[\underbrace{\bar{u}_2 \not\epsilon(\lambda)\gamma^5 u_1}_{\mathcal{S}_\lambda^{s_1 s_2}}\underbrace{\int\Big(2\sum_{\alpha=1}^{7}T_\alpha+\sum_{\alpha=8}^{14}T_\alpha\Big)}_{\mathcal{I}} - \underbrace{\bar{u}_2\not\epsilon(\lambda)\frac{\not\Delta_T}{M}\gamma^5 u_1}_{\mathcal{S'}_\lambda^{s_1 s_2}}\underbrace{\int\Big(2\sum_{\alpha=1}^{7}T'_\alpha+\sum_{\alpha=8}^{14}T'_\alpha\Big)}_{\mathcal{I}'}\Big], \quad (9)$$

where $\int \equiv \int\limits_{-1+\xi}^{1+\xi} dx_1 dx_2 dx_3 \delta(2\xi - x_1 - x_2 - x_3)\int\limits_0^1 dy_1 dy_2 dy_3 \delta(1-y_1-y_2-y_3)$, $u(p_1,s_1) \equiv u_1$, $\bar{u}(p_2,s_2) \equiv \bar{u}_2$ and the coefficients T_α and $T'_\alpha (\alpha = 1, ..., 14)$ are functions of x_i, y_j, ξ and Δ^2 and are given in Table 1 of [4].

The expression in Eq. (9) is to be compared with the leading-twist amplitude for the baryonic-form factor [16]

$$\mathcal{M}_\lambda \propto -i(\bar{u}_2\not\epsilon(\lambda)u_1)\frac{\alpha_s^2 f_N^2}{Q^4}\int\Big(2\sum_{\alpha=1}^{7}T_\alpha^p(x_i,y_j,\xi,t)+\sum_{\alpha=8}^{14}T_\alpha^p(x_i,y_j,\xi,t)\Big). \quad (10)$$

The factors T_α^p are very similar to the T_α obtained here. However, the integration domain is different. In the form factor case

$$\int \quad \text{stands for} \quad \int\limits_0^1 dx_1 dx_2 dx_3 \delta(1-\sum_i x_i)\int\limits_0^1 dy_1 dy_2 dy_3 \delta(1-\sum_i y_i). \quad (11)$$

Consequently, the integration of denominators in T_α^p such as $1/(x_i + i\varepsilon)$ do not generate any imaginary parts. On the contrary, the integrations of similar denominators in T_α and T'_α over the TDA integration domain will generate an imaginary part when passing from the ERBL region (all $x_i > 0$) to one of the DGLAP regions (one $x_i < 0$). This will be the source of the SSA as we will show later on.

3.3. The Single Transverse Spin Asymmetry

Since we are interested in the leading twist contribution of this asymmetry, we can sum only over the transverse polarisation of the virtual photon using $\sum_{\lambda=x,y}\varepsilon(\lambda)^\mu(\varepsilon(\lambda)^\nu)^* = -g^{\mu\nu} + (p^\mu n^\nu + p^\nu n^\mu)/(p.n)$. The sum on the final-proton spin s_2 is done using $\sum_{s_2} u_\alpha(p_2,s_2)\bar{u}_\beta(p_2,s_2) = (\not p_2 + M)_{\alpha\beta}$. As regards the initial-proton spinor, one uses the following relation involving its transverse spin s_1, $u_\alpha(p_1,s_1)\bar{u}_\beta(p_1,s_1) = 1/2(1+\gamma^5\not s_1)(\not p_1 + M)_{\alpha\beta}$.

Dropping the contributions proportional to the proton mass, the spin asymmetry reads

$$\sigma^{s_1} - \sigma^{-s_1} = 8\frac{|C|^2}{Q^6}\frac{1+\xi}{\xi}\frac{\epsilon^{n p s_1 \Delta_T}}{M}\Im m(\mathcal{I}'\mathcal{I}^*) \quad (12)$$

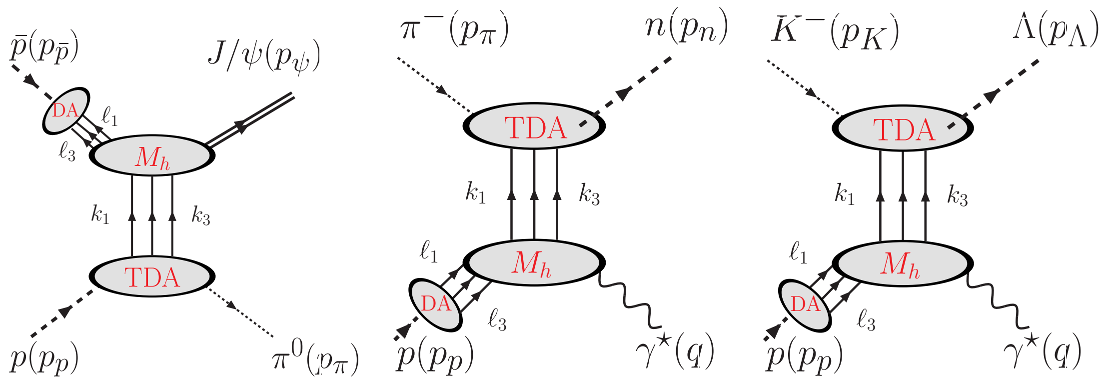
$$= -4\frac{|C|^2}{Q^6}\frac{|\Delta_T|}{M}\frac{1+\xi}{\xi}\sin(\phi-\phi_S)\Im m(\mathcal{I}'\mathcal{I}^*). \quad (13)$$

Comparing with the expressions for the unpolarised cross section obtained in [4], one concludes that the asymmetry for the hard-parton induced contribution is leading-twist as soon as $\Delta_T \neq 0$ and $\mathcal{I}$ or $\mathcal{I}'$ are no longer pure real or pure imaginary numbers. This is precisely what one expects when DGLAP contributions are taken into account [17].

3.4. The Single Transverse Spin Asymmetry in $\bar{p}p$ annihilation

Following the same reasonning, one can show that such an asymmetry would arise in $\bar{p}p \to \ell^+\ell^-\pi^0$. Let us however emphasise that the asymmetry on the (target) proton spin would then be leading twist only for the small-t regime, while its projectile spin asymmetry is leading twist in the small-u regime. In the case where only the target can be polarised, SSA studies should then preferentially be carried out by looking at slowly moving pions associated with a dilepton.

In general, spin asymmetries related to TDAs studies are of interest for the spin of the baryon which undergoes a non-perturbative transition into a meson.



(a) $p\bar{p} \to J/\psi\pi^0$ at small $t = (p_p - p_{\pi^0})^2$ (b) $\pi^- p \to \gamma^\star n$ at small $u = (p_\pi - p_n)^2$ (c) $K^- p \to \gamma^\star \Lambda$ at small $u = (p_K - p_\Lambda)^2$

Figure 3. Three additional hard-exclusive reactions involving the TDAs at GSI-FAIR and COMPASS.

4. Other processes involving the TDAs

4.1. $\bar{p}p \to J/\psi\pi^0$ at GSI-FAIR

It is well-known that exclusive decay of J/ψ into $p\bar{p}$ is rather well accounted for by the pQCD mechanism involving DAs and a hard scattering with three gluons exchange between the $c\bar{c}$ pair and the $(q\bar{q})(q\bar{q})(q\bar{q})$ system [16, 19]. One then naturally expects that the production of J/ψ in association with a pion could also be described using the DAs and the TDAs on the one hand and the same hard scattering on the other. This is schematically illustrated on Fig. 3 (a). Scaling properties and specific polarisation dominance should also appear as well as SSA if the reaction is indeed occuring at the parton level and not at the level of the baryon and the meson[18].

4.2. $\pi^- p \to n\gamma^\star$ and $K^- p \to \Lambda\gamma^\star$ at COMPASS

Meson beams can also be used to study hard exclusive reaction in the timelike region. Due to baryon-number conservation, one should still have a baryon in the final state. Using a π^-, it was proposed [20] to study $\pi^- p \to n\gamma^\star$ in the forward region –small $t = (p_p - p_n)^2$– to extract information about the GPDs. Accordingly, it is perfectly reasonable to think of the same study in the backward regime, where the final state baryon is very energetic in the target-rest frame, more specifically at small $u = (p_\pi - p_n)^2$. For an outgoing photon with sufficiently large Q^2, the amplitude should factorise in terms of pion-to-neutron TDAs as depicted on Fig. 3 (b).

As for $\bar{p}p \to \gamma^\star\pi^0$, the outgoing photon would be transversely polarised. However, the target transverse spin asymmetry would be higher twist since it is the neutron which undergoes the non-perturbative transition to the π^0. Only the asymmetry of the outgoing neutron spin would not be suppressed by any power of Q^2. To avoid this shortcoming, one could think of using K^- beam. The corresponding process would then be $K^- p \to \Lambda\gamma^\star$, with a fast Λ in the final state (see Fig. 3 (c)). Spin asymmetry on the latter could be then analysed through the azimuthal angular dependence of the decay $\Lambda \to p\pi^-$.

5. Discussion and conclusion

Although the knowledge of *baryon to meson* TDAs has recently improved significantly thanks to a first study in the meson cloud model [9] and another one focused on their spectral representations[10], model-independent observables aimed at studying the backward regime of meson electroproduction or exclusive $\bar{p}p$ annhiliation into a meson and a dilepton will still be the bread-and-butter of this field for the months to come.

Two obvious model-independent observables are the fixed Q^{-6} dependence of the amplitude squared, for which the transverse polarisation of the $\gamma^\star$ dominates, as well as the scaling in ξ, to be compared to x_{Bj} in DIS. We also find it particularly relevant to emphasise that the study of the asymmetry of the target transverse spin would reveal unique information on the nature of the particles exchanged in the u channel, be it a "mere" baryon slightly off-shell, or three perturbative quarks. For non-vanishing transverse momenta (Δ_T), one expects in the latter case an asymmetry of the same order as the unpolarised cross section, while, in the former case, they would be most likely decreasing for increasing W^2 and Q^2.

In conclusion, let us stress that we expect spin observables to be a major tool to establish the dominance of a partonic process in the hard-exclusive reactions on which we have focused here. This reminds us of their importance in opening the field of Generalised Parton Distribution studies. Scaling tests – although crucial for the proof of the adequacy of the theoretical description – are often experimentally undecisive because of the small lever arm reachable in pionneering experiments. Even in the case of unpolarised beam and targets, the dominance of transversely-polarised virtual-photon exchange may be checked through azymuthal dependence studies. As emphasised here, using transversely-polarised targets or analysing the polarisation of final-states particles will allow for other likely decisive tests of the QCD understanding of these reactions.

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