Backward DVCS and Proton to Photon Transition Distribution Amplitudes

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We analyse deeply-virtual Compton scattering on a proton target, $\gamma^* P \rightarrow P' \gamma$ in the backward region and in the scaling regime. We define the transition distribution amplitudes which describe the proton to photon transition. Model-independent predictions are given to test this description, for current or planned experiments at JLab or by Hermes.

1. INTRODUCTION

Deeply virtual Compton scattering (DVCS) at small momentum transfer $t$ has been the subject of a continuous progress in recent years, both on the theoretical side with the understanding of factorisation properties which allow a consistent calculation of the amplitude in the framework of QCD, and on the experimental side with the success of experiments at HERA and JLab. The generalised parton distributions (GPDs) which describe the soft part of the scattering amplitude indeed contain much information on the hadronic structure, which would remain hidden without this new opportunity [1]. In Ref. [2], it has been advocated that the same virtual Compton scattering reaction

$$eP(p_1) \rightarrow e'P'(p_2)\gamma(p_\gamma)$$

as well as electroproduction of meson ($\pi$, $\rho$, . . .)

$$eP(p_1) \rightarrow e'P'(p_2)M(p_M)$$

in the backward kinematics (namely small $u = (p_\gamma - p_1)^2$ or $u = (p_M - p_1)^2$) could be analysed in a slightly modified framework, the amplitude being factorised (see Fig. 1 (a) and (b)) at leading twist as

$$\mathcal{M}(Q^2, \xi, \Delta^2) \propto \int dx_i dy_j \Phi(y_j, Q^2)M_h(x_i, y_j, \xi) T(x_i, \xi, \Delta^2) ,$$

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where $\Phi(y_j, Q^2)$ is the proton distribution amplitude, $M_h$ is a perturbatively calculable hard scattering amplitude and $T(x_i, \xi, \Delta^2)$ are transition distribution amplitudes (TDAs) defined as the matrix elements of light-cone operators between a proton and a photon state or between a proton and a meson state.

The variable $x_i$ describes the fraction of light-cone momentum carried by the quark $i$ off the initial proton, $y_j$ is the corresponding one for the quark $j$ entering the final state proton, $\Delta = p_\gamma - p_1$ and the skewness variable $\xi$ describes the loss of plus-momentum of the incident proton (see section 2 for more details on kinematics).

In the large angle regime (around 90 degrees), the large value of $-t = -(p_1 - p_2)^2$ sets the perturbative scale. In the small angle regime as well as for the backward regime, it is the large virtuality $Q^2$ of the initial photon which allows a perturbative expansion of a subprocess scattering amplitude. Of course in the backward regime, small $-u = -(p_\gamma - p_1)^2$ means large $-t$, and even $-t$ larger than at 90 degrees, but this does not introduce a new scale in the problem, exactly as for the forward DVCS case for which, $-t$ being small, $-u$ is very large.

Figure 1. (a) Factorised amplitude for deeply-virtual Compton scattering on proton in the backward region; (b) Factorised amplitude for meson electroproduction on proton in the backward region. (c) Factorised amplitude for meson-pair ($\pi\pi$) production in $\gamma^*\gamma$ collisions.

In Ref. [3], we have defined the leading-twist proton to pion $P \rightarrow \pi$ transition distribution amplitudes from the Fourier transform$^4$ of the matrix element

$$\langle \pi | \epsilon_{ijk} q_\alpha^i(z_1 n) [z_1; z_0] q_\beta^j(z_2 n) [z_2; z_0] q_\gamma^k(z_3 n) [z_3; z_0] | P \rangle,$$

(4)

The brackets $[z_i; z_0]$ in Eq. (4) account for the insertion of a path-ordered gluonic exponential along the straight line connecting an arbitrary initial point $z_0 n$ and a final

$^4$In the following, we shall use the notation $\mathcal{F} \equiv (p.n)^3 \int_{-\infty}^{\infty} dz e^{\Sigma; x; z; p.n}$. 
one $z_in$:

$$[z_i; z_0] \equiv \exp \left[ ig \int_0^1 dt (z_i - z_0) n_\mu A^\mu (n[tz_i + (1 - t)z_0]) \right].$$

(5)

which provide the QCD-gauge invariance for non-local operator and equal unity in a light-like (axial) gauge.

In a similar way, we shall define in section 3 the proton to photon TDAs from the Fourier transform of the matrix element

$$\langle \gamma | \bar{q}_\alpha (z_1 n) [z_1; z_0] q_\beta (z_2 n) [z_2; z_0] q_\gamma (z_3 n) [z_3; z_0] | P \rangle.$$  

(6)

In the simpler mesonic case, a perturbative limit has been obtained \[4\] for the $\rho \to \gamma^*$ transition. For $\pi \to \gamma$ one, where there are only four leading-twist TDAs \[2\] entering the parametrisation of the matrix element

$$\langle \gamma | \bar{q}_\alpha (z_1 n) [z_1; z_0] q_\beta (z_0 n) | \pi \rangle,$$

we have recently shown \[5\] that experimental analysis of processes such as $\gamma^* \gamma \to \rho \pi$ and $\gamma^* \gamma \to \pi \pi$, see Fig. 1 (c), involving these TDAs could be carried out, e.g. the background from the Bremsstrahlung is small if not absent and rates are sizable at present $e^+e^-$ facilities.

2. Kinematics

The momenta of the processes $\gamma^* P \to P' \gamma$ are defined as shown in Fig. 1 (a). The $z$-axis is chosen along the initial nucleon momentum and the $x-z$ plane is identified with the collision plane. Then, we define the light-cone vectors $p$ and $n$ ($p^2 = n^2 = 0$) such that

$$2 p.n = 1,$$

and its transverse component $\Delta^T$, which we choose to be along the $x$-axis. From those, we define $\xi$ in an usual way as

$$\xi = -\frac{\Delta^T}{p.n}.$$

We can then express the momenta of the particles through their Sudakov decomposition:\footnote{$\Delta^T_2 < 0.$}

$$p_1 = (1 + \xi) p + \frac{M^2}{1 + \xi} n, \quad p_\gamma = (1 - \xi) p - \frac{\Delta^2}{1 - \xi} n + \Delta T,$$

$$p_2 = (2\xi - 1)p + n (Q^2 + \frac{\Delta^2}{1 + \xi} - \frac{M^2}{1 + \xi}) - \Delta T, \quad q = -p + Q^2 n.$$  

(7)

Using the natural gauge choice $\varepsilon.n = 0$, the photon polarisation vector $\varepsilon(p_\gamma)$ can be chosen to be either a normalised vector along the $y$-axis,

$$\varepsilon_{T1} = \varepsilon_y \quad \text{or} \quad \varepsilon_{T2} = \frac{\Delta T}{\sqrt{-\Delta^2_T}} + 2 \sqrt{-\Delta^2_T} \frac{\Delta^T}{1 - \xi} n,$$

(8)

which gives $\varepsilon_{T2} = \varepsilon_x$ at $\Delta_T = 0$.

In an arbitrary QED gauge, where $\varepsilon' = \varepsilon_T + \lambda p_\gamma$, we have at $\Delta_T = 0$

$$\varepsilon'_1 = \lambda (1 - \xi) p + \varepsilon_y, \quad \varepsilon'_2 = \lambda (1 - \xi) p + \varepsilon_x.$$  

(9)

Therefore one has, in any gauge and at $\Delta_T = 0$,

$$\varepsilon.p = 0, \quad \varepsilon.n = \lambda \frac{1 - \xi}{2}, \quad \varepsilon.\Delta_T = 0.$$  

(10)
3. The Proton to Photon TDAs

The spinorial and Lorentz decomposition of the matrix element will follow the same line as the one for $P \to \pi$ TDA \cite{3} and for baryon DA \cite{6}. The fractions of plus momenta are labelled $x_1$, $x_2$ and $x_3$, and their supports are within $[-1 + \xi, 1 + \xi]$. Momentum conservation implies (we restrict to the case $\xi > 0$):

$$\sum_i x_i = 2\xi. \quad (11)$$

The configurations with positive momentum fractions, $x_i \geq 0$, describe the creation of quarks, whereas those with negative momentum fractions, $x_i \leq 0$, the absorption of antiquarks.

Counting the degrees of freedom fixes the number of independent $P \to \gamma$ TDAs to 16, since each quark, photon and proton have two helicity states (leading to $2^5$ helicity amplitudes) and parity relates amplitudes with opposite helicities for all particles. We can equally say that the photon has spin 1, which would normally give 24 TDAs as in the $P \to V$ where $V$ is a massive vector particle, but gauge invariance provides us with 8 relations between TDAs, which reduces again the number to 16.

The case $\Delta_T = 0$ is simpler since the matrix elements can be written only in terms of 4 TDAs. Indeed, since at $\Delta_T = 0$, there is no angular momentum exchanged, the helicity is conserved. We have three possible processes as $P(\uparrow) \to uud(\uparrow\downarrow\downarrow) + \gamma(\uparrow)$ where the quark with helicity -1 is either the $u$’s or the $d$, but also $P(\uparrow) \to uud(\uparrow\uparrow\uparrow) + \gamma(\downarrow)$. Therefore taking this limit on the complete set of the 16 TDAs should reduce it to 4.

In order to build leading-twist structures (maximising the power of $P^+$), we have first to separate the spinor $N(p_1)$ in its small ($N^- \sim \sqrt{1/P^+}$) and large ($N^+ \sim \sqrt{P^+}$) component:

$$N = (\psi\gamma^\mu + \gamma^\mu\psi)N = N^- + N^+. \quad (12)$$

Using the Dirac equation $\gamma_\mu N(p_1) = MN(p_1)$ and Eq. (17), it is easy to see that

$$\gamma^\mu N = \frac{M}{2(1 + \xi)}N^+ + O(1/P^+) \quad \text{and} \quad \gamma^\mu N = \frac{1 + \xi}{2M}N^- + O(1/P^+). \quad (13)$$

We then proceed in the following way:

1. the structures are to be linear in the photon polarisation vector (through scalar products with the momenta $(n, p$ and $\Delta_T)$, $\gamma^\mu$ or $\sigma^{\mu\nu}$).

2. we force the presence of $p$ ($\simeq P$) to help the twist counting in powers of $P^+$ (therefore the different leading-twist structures will scale like $(P^+)^{3/2}$);

3. $p$ does not appear in $\gamma^\mu N$ since this would remove one power of $P^+$;

4. $p$ does not appear in any scalar products $p.n$, $p.\Delta_T$ and $p.\varepsilon$ which would also destroy one power of $P^+$;

5. $p$ then only appears inside the parenthesis $(\cdot)_{\alpha\beta}$;
6. we impose the independence of the factors in \((\cdot)_{\alpha\beta}\) from two different structures; this can be checked by taking the trace of the product of two structures, and is therefore insured by choosing only independent Fierz (or Dirac) structures \(\gamma^5, \gamma^\mu, \gamma^5\gamma^\mu, \sigma^{\mu\nu}\).

7. Finally, to what concerns the spinor, it has only two large components. Hence, after a given \((\cdot)_{\alpha\beta}\), it appears only twice with a different Dirac structure (e.g. \(N\) and \(\not{\nu}\)).

This construction leads to define 24 possible independent structures for the transition proton to vector (whose factors \(V_i, A_i\) and \(T_i\) are dimensionless and real function of the momentum fractions \(x_i, \xi\) and \(\Delta^2\)).

\[
4\mathcal{F}\left(\langle V(p_T)\mid \epsilon_{ijk}u^i_\alpha(z_1n)u^j_\beta(z_2n)d^k_\gamma(z_3n)\rangle P(p_1, s_1)\right) = M \times \nonumber
\]

\[
\left( V^1_\xi(\delta\mathcal{P})_{\alpha\beta}(\xi N^+) + M^{-1}V^T_1(\epsilon.\Delta T)(\delta\mathcal{P})_{\alpha\beta}(N^+) + MV^1_n(\epsilon.n)(\delta\mathcal{P})_{\alpha\beta}(N^+) + \right. \\
M^{-1}V^2_\xi(\delta\mathcal{P})_{\alpha\beta}(\sigma^{\Delta\xi}N^+) + M^{-2}V^2_T(\epsilon.\Delta T)(\delta\mathcal{P})_{\alpha\beta}(\Delta\xi N^+) + V^2_n(\epsilon.n)(\delta\mathcal{P})_{\alpha\beta}(\Delta\xi N^+) + \\
A^\xi_1(\delta\gamma^5\mathcal{P})_{\alpha\beta}(\gamma^5\xi N^+) + M^{-1}A^T_1(\epsilon.\Delta T)(\delta\gamma^5\mathcal{P})_{\alpha\beta}(\gamma^5 N^+) + M A^n_1(\epsilon.n)(\delta\gamma^5\mathcal{P})_{\alpha\beta}(\gamma^5 N^+) + \\
M^{-1}A^2_2(\delta\gamma^5\mathcal{P})_{\alpha\beta}(\gamma^5\sigma^{\Delta\xi}N^+) + M^{-2}A^2_T(\epsilon.\Delta T)(\delta\gamma^5\mathcal{P})_{\alpha\beta}(\gamma^5\Delta\xi N^+) + \\
A^n_2(\epsilon.n)(\delta\gamma^5\mathcal{P})_{\alpha\beta}(\gamma^5\Delta\xi N^+) + T^1_1(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\sigma^{\mu\nu}N^+) + M^{-1}T^T_1(\epsilon.\Delta T)(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\gamma^\mu N^+) + \\
MT^n_1(\epsilon.n)(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\gamma^\mu N^+) + T^2_2(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(N^+) + M^{-2}T^T_2(\epsilon.\Delta T)(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\sigma^{\Delta\xi}N^+) + \\
T^2_n(\epsilon.n)(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\sigma^{\mu\nu}N^+) + M^{-1}T^T_3(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\Delta\xi N^+) + \\
M^{-2}T^T_3(\epsilon.\Delta T)(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\gamma^\mu N^+) + T^n_3(\epsilon.n)(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\gamma^\mu N^+) + M^{-1}T^T_4(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\Delta\xi N^+) + \\
M^{-3}T^T_4(\epsilon.\Delta T)(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\Delta\xi N^+) + M^{-1}T^T_4(\epsilon.n)(\sigma_{\mu\nu}\mathcal{P})_{\alpha\beta}(\Delta\xi N^+),
\]

where \(\sigma^{\mu\nu} \equiv \frac{1}{2}[\gamma^\mu, \gamma^\nu]\) and \(C\) is the charge-conjugation matrix.

The quark fields in the matrix element of Eq. \((\text{14})\) are defined according to the prescription of Mandelstam \([7]\) in order to make the latter QED gauge invariant. In the proton to photon case, gauge invariance of the r.h.s of Eq. \((\text{14})\) implies that the latter vanishes when \(\epsilon(p_\gamma)\) is replaced by \(p_\gamma\).

At the leading-twist accuracy, this provides us with 8 relations:

\[
V^\xi_1(1 - \frac{\xi}{\Delta^2}) M = V^T_1 \frac{\Delta^2}{M} + V^n_1 \frac{(1 - \xi)M}{2(1 + \xi)} = V^\xi_1 + V^T_1 \frac{\Delta^2}{M} + V^n_1 \frac{1 - \xi}{2} = 0, \\
A^\xi_1(1 - \frac{\xi}{\Delta^2}) M = A^T_1 \frac{\Delta^2}{M} + A^n_1 \frac{(1 - \xi)M}{2(1 + \xi)} = A^\xi_1 + A^T_1 \frac{\Delta^2}{M} + A^n_1 \frac{1 - \xi}{2} = 0, \\
T^\xi_1(1 - \frac{\xi}{\Delta^2}) M = T^T_1 \frac{\Delta^2}{M} + T^n_1 \frac{(1 - \xi)M}{2(1 + \xi)} = T^\xi_1 + T^T_1 \frac{\Delta^2}{M} + T^n_1 \frac{1 - \xi}{2} = 0.
\]

This effectively reduces the number of \(P \rightarrow \gamma\) TDAs to 16 as expected from the number of helicity amplitudes for the process \(P \rightarrow qqq\gamma\), and we have

\[
4\mathcal{F}\left(\langle \gamma(p_\gamma)\mid \epsilon_{ijk}u^i_\alpha(z_1n)u^j_\beta(z_2n)d^k_\gamma(z_3n)\rangle P(p_1, s_1)\right) = M \times \nonumber
\]

\[
\left( V^\xi_1(x_i, \xi, \Delta^2)(\delta\mathcal{P})_{\alpha\beta}(\xi N^+) + \frac{M}{1 + \xi}(\epsilon.n)(N^+) - \frac{2(\epsilon.n)}{1 - \xi}(\Delta\xi N^+)\right) + \nonumber
\]
\begin{align}
V_i^T(x_i, \xi, \Delta^2) \left[ \epsilon \frac{2\Delta^2}{1-\xi} (\epsilon .n) \right] (\hat{p} C)_{\alpha \beta} (N^+) \gamma + \\
V_i^T(x_i, \xi, \Delta^2) \left[ \epsilon \frac{2\Delta^2}{1-\xi} (\epsilon .n) \right] (\hat{p} C)_{\alpha \beta} (\Delta \gamma N^+) \gamma + \\
V_i^T(x_i, \xi, \Delta^2) \left[ \epsilon \frac{2\Delta^2}{1-\xi} (\epsilon .n) \right] (\hat{p} C)_{\alpha \beta} (\Delta \gamma N^+) \gamma + \\
A_i^T(x_i, \xi, \Delta^2) (\Gamma^5 C)_{\alpha \beta} \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M (\epsilon .n)}{2(1+\xi)} (\gamma^5 \Delta \gamma N^+) \gamma \right] + \\
A_i^T(x_i, \xi, \Delta^2) (\Gamma^5 C)_{\alpha \beta} \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M (\epsilon .n)}{2(1+\xi)} (\gamma^5 \Delta \gamma N^+) \gamma \right] + \\
A_i^T(x_i, \xi, \Delta^2) (\Gamma^5 C)_{\alpha \beta} \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M (\epsilon .n)}{2(1+\xi)} (\gamma^5 \Delta \gamma N^+) \gamma \right] + \\
T_i^T(x_i, \xi, \Delta^2) (\sigma_{\alpha \beta} C)_{\alpha \beta} \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M (\epsilon .n)}{2(1+\xi)} (\gamma^5 \Delta \gamma N^+) \gamma \right] + \\
T_i^T(x_i, \xi, \Delta^2) (\sigma_{\alpha \beta} C)_{\alpha \beta} \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M (\epsilon .n)}{2(1+\xi)} (\gamma^5 \Delta \gamma N^+) \gamma \right] + \\
T_i^T(x_i, \xi, \Delta^2) (\sigma_{\alpha \beta} C)_{\alpha \beta} \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M (\epsilon .n)}{2(1+\xi)} (\gamma^5 \Delta \gamma N^+) \gamma \right] + \\
T_i^T(x_i, \xi, \Delta^2) (\sigma_{\alpha \beta} C)_{\alpha \beta} \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M (\epsilon .n)}{2(1+\xi)} (\gamma^5 \Delta \gamma N^+) \gamma \right].
\end{align}

As discussed earlier, the $\Delta_T = 0$ case is much simpler since it involves only 4 TDAs to describe the proton to photon transition. Moreover, in the Bjorken scaling which interests us, $\Delta_T$ is in any case supposed to be small, making this limit $\Delta_T = 0$ particularly fruitful to consider.

The four expected TDAs for $p \to \gamma$ TDAs are straightforwardly obtained from Eq. (16) by setting $\Delta_T = 0$:

\begin{align}
4 \mathcal{F} \left[ (\gamma(p)) | \epsilon_{ij} u^i_\alpha (z_1 n) u^j_\beta (z_2 n) d^k_\chi (z_3 n) | P(p_1, s_1) \right] = M \times \\
(V_1^T(x_i, \xi, \Delta^2) (\hat{p} C)_{\alpha \beta} \left[ (\hat{p} N^+) \gamma - \frac{M}{1+\xi} (\epsilon .n) (N^+) \gamma \right] + \\
A_1^T(x_i, \xi, \Delta^2) (\Gamma^5 C)_{\alpha \beta} \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M}{1+\xi} (\epsilon .n) (\gamma^5 N^+) \gamma \right] + \\
A_2^T(x_i, \xi, \Delta^2) \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M}{1+\xi} (\epsilon .n) (\gamma^5 N^+) \gamma \right] + \\
A_3^T(x_i, \xi, \Delta^2) \left[ (\gamma^5 \hat{p} N^+) \gamma - \frac{M}{1+\xi} (\epsilon .n) (\gamma^5 N^+) \gamma \right].
\end{align}
\[ T_1^\varepsilon(x_i, \xi, \Delta^2) (\sigma_{\mu\nu} C)_{\alpha\beta} \left[ (\sigma^{\mu\nu} N^+)_{\gamma} - \frac{M}{1 + \xi} \frac{\varepsilon.n}{2} (\gamma^\mu N^+)_\gamma \right] + T_5^\varepsilon(x_i, \xi, t) (\sigma_{\mu\nu} C)_{\alpha\beta}(N^+)_\gamma. \]

4. Amplitude calculation at $\Delta_T = 0$ and model-independent predictions

Let us now consider the calculation of the helicity amplitude in the $\Delta_T = 0$ limit. At leading order in $\alpha_S$, the helicity amplitude $M_{\lambda_1,\lambda_2,s_1,s_2}$ for the reaction

\[ \gamma^* (q, \lambda_1) P(p_1, s_1) \rightarrow P(p_2, s_2) \gamma(p_\gamma, \lambda_2) \]

is calculated similarly to the baryonic form-factor [11, 12]. It reads

\[ M_{\lambda_1,\lambda_2,s_1,s_2} \propto e\bar{u}(p_2, s_2)\gamma^\lambda_2 \gamma^5 u(p_1, s_1) M(\alpha_S(Q^2))^2 Q^4 \times \int_{1+\xi}^{1} d^3x_i \int_0^1 d^3y_j \delta(\sum x_i - 2\xi) \delta(\sum y_j - 1) \sum_{\alpha=1}^{10} T_\alpha(x_i, y_j, \xi, \Delta^2) \]

where the coefficients $T_\alpha$ include both the proton to photon TDAs and the final-state proton DAs. The structure $e\bar{u}(p_2, s_2)\gamma^\lambda_2 \gamma^5 u(p_1, s_1) \gamma_5$ selects opposite helicity states for the final and initial protons. The same statement holds for the photons. This is a model independent result at $\Delta_T = 0$ as well as the scaling $(\alpha_S(Q^2))^2 Q^4$ up to logarithmic corrections due to the evolution of the TDAs and DAs.

5. Conclusions and perspectives

We have defined the 16 proton to photon Transition Distribution Amplitudes entering the description of backward virtual Compton scattering on proton target. Since the study in terms of GPDs of the latter process in the forward region has been very fruitful to understand the underlying structure of the hadron, we foresee that the corresponding one with TDAs of the backward region be of equal importance, if not more since it involves the exchange of 3 quarks.

We have also calculated the amplitude for the process $\gamma^*P \rightarrow P'\gamma$ in terms of the TDAs. In order to provide with theoretical evaluations of cross sections, we still have to develop an adequate model for the TDAs $V_i$, $A_i$ and $T_i$. This may be done through the introduction of quadruple distributions, which generalise the double distributions introduced by Radyushkin [13] in the GPD case. Similarly to this latter case, it will also ensure the proper polynomiality and support properties of the TDAs. A limiting value of the TDA for $\xi \rightarrow 1$ may be derived by considering the soft photon limit of the scattering amplitude and may be used as a model input in these quadruple distributions, whereas for the GPDs the diagonal limit, i.e. the parton distribution functions, was used as input.

Model independent predictions follow from the way we propose to factorise the amplitude: only helicity amplitudes with opposite signs for both protons and photons will

\^{6}whose details (omitted due to lack of space) will be presented in a forthcoming publication [10].
be nonzero at $\Delta_T = 0$. Furthermore, the amplitude scales as $\frac{(\alpha_S(Q^2))^2}{Q^4}$ as do the similar amplitudes for backward electroproduction of mesons, i.e. $\gamma^* P \to P'\pi$ or $\gamma^* P \to P'\rho$. Observation of such a universal scaling law would provide with indications that the picture holds and dominates over a purely hadronic model as considered in [9], where data are presented for low energies but for $Q^2 = 1\text{GeV}^2$.

Let us finally stress that, first, in the backward region considered here, there is almost no Bethe-Heitler contribution: the experimentally measured cross sections will depend bilinearly on the TDAs; secondly, the same matrix elements of Eq. (14) factorise in the amplitude $PP \to \gamma^*\gamma$, which may be studied at GSI. This universality makes the TDAs an essential tool for the generalisation of the hadronic studies carried at electron machines to complementary studies to be carried in proton antiproton experiments. Thus, the properties of these TDAs are planned to be studied both at the upgraded JLab experiments and with PANDA and PAX [13] at GSI.

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