Timelike Compton Scattering off the Proton: beam and/or target spin asymmetries.

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We present a sample of results of our work to be published soon on Timelike Compton scattering off the proton, in the framework of the Generalized Parton Distributions formalism.

1 Introduction

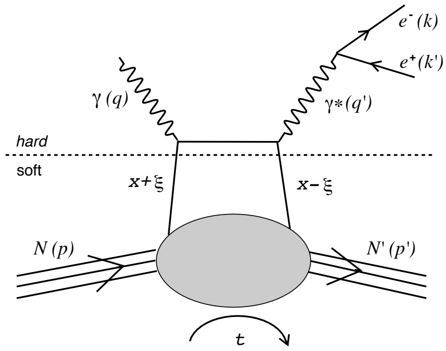

Figure 1: Leading twist TCS diagram.

More than 40 years after the discovery of point-like components within the proton, its quarks and gluons structure is still not well understood and is still intensively studied. Hard exclusive processes on the proton provide access to the Generalized Parton Distributions (GPDs) [1, 2, 3, 4] which contain informations about the longitudinal momentum and the spatial transverse distributions of partons inside the proton (in a frame where the nucleon has an "infinite" momentum along its longitudinal direction). Such a hard exclusive process is the Deeply Virtual Compton scattering process which corresponds to the reaction $\gamma^{(*)}P \to \gamma^{(*)}P$ and to the scattering of a high-energy virtual photon off a quark inside the proton. There are two particular cases of deep Compton processes. "Spacelike" Deeply Virtual Compton Scattering (DVCS) corresponds to the case where the incoming photon is emitted by a lepton beam and has a high spacelike virtuality and and where the final photon is real. The DVCS process has been studied for the past ∼15 years and is still intensively studied both theoretically and experimentally. The second particular case of deep Compton scattering is the Timelike Compton Scattering (TCS) process. It corresponds to the case where the incoming photon is real and the final photon has a high timelike virtuality and decays into a lepton pair (see Fig. 1). Contrary to DVCS, there is no published experimental data yet for TCS. Both DVCS and TCS give access to the same proton GPDs in the QCD leading twist formalism. The study of TCS in parallel to DVCS is a very powerful way to check the universality of GPDs and/or to study higher twist effects. The reaction $\gamma P \to e^+e^-P$ also involves the Bethe-Heitler process, where the incoming real photon creates a lepton pair, which then interacts with the proton. It is not sensitive to the GPDs but to the form factors. It can be calculated with a few percent accuracy.

2 Amplitudes and observables

The four vectors involved are indicated in Fig. 1. According to QCD factorization theorems, at sufficiently large $Q'^2 = (k + k')^2$ (photon's virtuality), we can decompose the TCS amplitude into a soft part, parameterized by the GPDs, and a hard part, exactly calculable by Feynman diagrams techniques. We work in a frame where the average protons and the average photons momenta, respectively P and $\bar{q}$, are collinear along the z-axis and in opposite directions. We define the lightlike vectors along the positive and negative z directions as $\tilde{p}^\mu = P^+/\sqrt{2}(1,0,0,1)$ and $n^\mu = 1/P^+ \cdot 1/\sqrt{2}(1,0,0,-1)$, with $P^+ \equiv (P^0 + P^3)/\sqrt{2}$. We have the properties $\tilde{p}^2 = n^2 = 0$ and $\tilde{p} \cdot n = 1$. In this frame, the TCS amplitude can be written in the asymptotic limit (mass terms are neglected with respect to Q'^2) with the Ji convention for GPDs [5]:

$$T^{TCS} = -\frac{e^3}{q'^2}\,\bar{u}(k)\,\gamma^\nu\,v(k')\,\epsilon^\mu(q)$$

$$\left[\frac{1}{2}\,(-g_{\mu\nu})_\perp \int_{-1}^{1} dx\,\left(\frac{1}{x-\xi-i\epsilon} + \frac{1}{x+\xi+i\epsilon}\right) \cdot \left(H(x,\xi,t)\bar{u}(p')\,\not{n}u(p) + E(x,\xi,t)\bar{u}(p')i\sigma^{\alpha\beta}n_\alpha\frac{\Delta_\beta}{2M}\,u(p)\right)\right.$$

$$\left.-\frac{i}{2}(\epsilon_{\nu\mu})_\perp \int_{-1}^{1} dx\,\left(\frac{1}{x-\xi-i\epsilon} - \frac{1}{x+\xi+i\epsilon}\right) \cdot \left(\tilde{H}(x,\xi,t)\bar{u}(p')\,\not{n}\gamma_5\,u(p) + \tilde{E}(x,\xi,t)\bar{u}(p')\gamma_5\frac{\Delta.n}{2M}\,u(p)\right)\right] \quad (1)$$

where x is the quark longitudinal momentum fraction, $\Delta = (p' - p)$ is the momentum transfer, $t = \Delta^2$ and ξ is defined as

$$\xi = -\frac{(p-p').(q'+q)}{(p+p').(q'+q)}. \quad (2)$$

In Eq. 1, we used the metric

$$(-g_{\mu\nu})_\perp = -g_{\mu\nu} + \tilde{p}_\mu n_\nu + \tilde{p}_\nu n_\mu \;, \quad (\epsilon_{\nu\mu})_\perp = \epsilon_{\nu\mu\alpha\beta}\,n^\alpha\,\tilde{p}^\beta. \quad (3)$$

The Bethe-Heitler amplitude reads:

$$T^{BH} = -\frac{e^3}{\Delta^2}\bar{u}(p')\left(\gamma^\nu F_1(t) + \frac{i\sigma^{\nu\rho}\Delta_\rho}{2\,M}F_2(t)\right)u(p)\,\epsilon^\mu(q)\bar{u}(k)\left(\gamma_\mu\frac{\not{k}-\not{q}}{(k-q)^2}\gamma_\nu + \gamma_\nu\frac{\not{q}-\not{k}'}{(q-k')^2}\gamma_\mu\right)v(k'), \quad (4)$$

where $F_1(t)$ and $F_2(t)$ are the proton Dirac and Pauli form factors. At fixed beam energy, the cross section of the photoproduction process depends on four independant kinematic variables, which we choose as: Q'^2, t and the two angles θ and ϕ of the decay electron in the γ^* center of mass. The 4-differential unpolarized cross section reads:

$$\frac{d^4\sigma}{dQ'^2dtd\Omega}(\gamma p \to p'e^+e^-) = \frac{1}{(2\pi)^4}\frac{1}{64}\frac{1}{(2ME_\gamma)^2}\mid T^{BH} + T^{TCS}\mid^2, \quad (5)$$

where $\mid T^{BH} + T^{TCS}\mid^2$ is averaged over the target proton and beam polarizations and summed over the final proton spins.

We define the single and double spin asymmetries as:

$$A_{\odot U}\;(A_{Ui}) = \frac{\sigma^+ - \sigma^-}{\sigma^+ + \sigma^-}, \quad A_{\odot i} = \frac{(\sigma^{++} + \sigma^{--}) - (\sigma^{+-} + \sigma^{-+})}{\sigma^{++} + \sigma^{--} + \sigma^{+-} + \sigma^{-+}}, \quad (6)$$

where the first index of A corresponds to the polarization state of the beam and the second one corresponds to the polarization state of the target. $A_{\odot U}$ is the circularly polarized beam spin asymmetry. The + and − superscripts in σ correspond to the two photon spin states, right and left polarized. A_{Ui} are the single target spin asymmetries where the + and − superscripts refer to the target spin orientations along the axis $i = x, y, z$. The axis x and y are perpendicular to the incoming proton direction (along the z-axis) in the γP center of mass frame and are respectivelly in the scattering plane and perpendicular to this plane. $A_{\odot i}$ are the double spin asymmetries with a circularly polarized beam and with a polarized target. We finally define the single linearly polarized beam spin asymmetry as

$$A_{\ell U}(\Psi) = \frac{\sigma_x(\Psi) - \sigma_y(\Psi)}{\sigma_x(\Psi) + \sigma_y(\Psi)}, \tag{7}$$

where Ψ is the angle between the photon polarization vector and the $\gamma P \to \gamma^* P'$ plane and where σ_x (σ_y) indicate a photon polarized in the x-(y-)direction.

3 Numerical results and sensitivity to GPDs

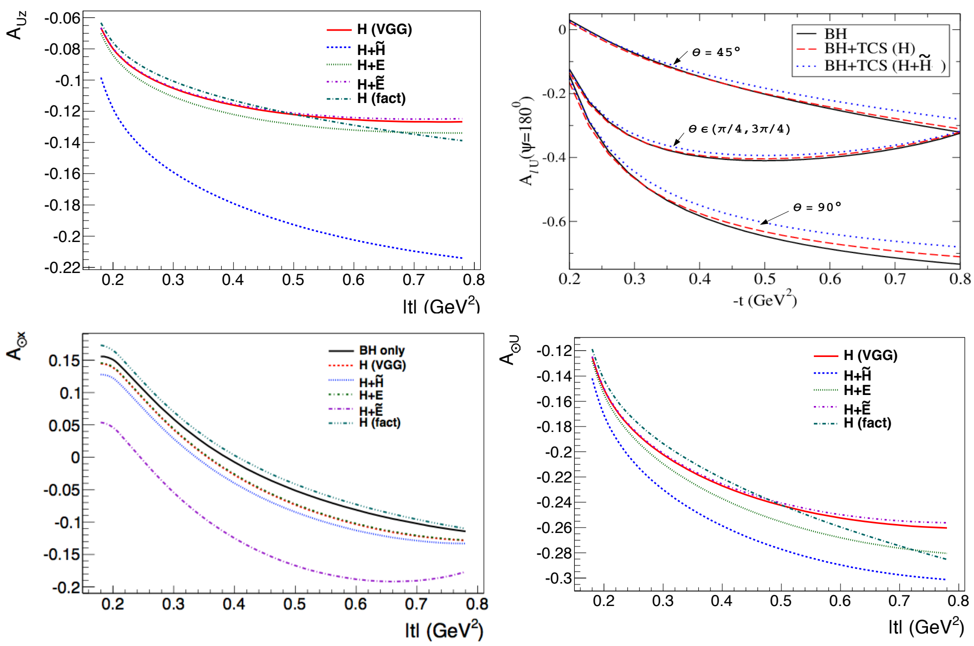

Figure 2: Spin asymmetries as a function of $-t$. Top left: $A_{\odot U}$ for BH+TCS. Top right: $A_{\ell U}$ for BH and BH+TCS. Bottom left: A_{Uz} for BH+TCS. Bottom right: $A_{\odot x}$ for BH and BH+TCS. All calculations are done at $\xi = 0.2$, $Q'^2 = 7$ GeV^2, $\phi = 90°$ and θ integrated over $[45°, 135°]$. $A_{\ell U}$ is also shown for $\theta = 45°$ and $\theta = 90°$.

We performed our calculations using the GPD parameterization of the VGG model [6, 7, 8]. We focus here on the spin asymmetries. Figure 2 shows the circularly (top row left) and linearly (top row right) polarized beam spin asymmetries as a function of t. One should note that $A_{\odot U}$

is particularly sensitive to the GPDs as it is exactly 0 for BH alone. It comes from the fact that this asymmetry is sensitive to the imaginary part of the amplitudes and the BH amplitude is purely real. We also show $A_{\odot U}$ with a factorized-t ansatz instead of a Reggeized-t ansatz for the H GPD which illustrates the sensitivity to the GPD modeling. In contrast, the $A_{\ell U}$ asymmetry, which is strong, is dominated by the BH and the TCS makes up only small deviations. Indeed, this asymmetry is sensitive to the real part of the amplitudes.

We display in Fig. 2 (bottom row) two examples of asymmetries with a polarized target: A_{Uz} and $A_{\odot x}$ (double spin asymmetry). We present the results for TCS+BH with different GPDs contributions and parameterizations. All single target spin asymmetries are zero for the BH alone as they are proportional to the imaginary part of the amplitudes. This makes the A_{Ui} asymmetries privileged observables to study GPDs. On the contrary, it is more difficult to access GPDs with double spin asymmetries as the BH alone produces a strong double spin asymmetry.

4 Discussion

We have presented a sample of our results to be published soon, namely the t-dependence of single and double spin asymmetries for the $\gamma P \to e^+e^-P$ reaction which we analyzed in the framework of the GPD formalism. We didn't discuss this here due to lack of space but we also compared our unpolarized cross sections and our single beam spin asymmetries with those of the earlier work of Refs [9, 10] and they are in agreement at the few percent level. We have introduced in our work the target polarization in order to define the single and double spin asymmetries with polarized targets. We have also introduced some higher twist corrections and gauge invariance restoration terms.

As the BH contribution alone doesn't contribute to single target spin asymmetries and to circularly polarized beam spin asymmetries, these observables are good candidates to study GPDs. Such measurements can be envisaged at the JLab 12 GeV facility. In particular, a proposal has been accepted for the CLAS12 experiment (JLab Hall B) to measure the unpolarized BH+TCS cross section [11]. The work that we presented here can open the way to new complementary experimental programs with polarized beams and/or targets.

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