

PROBING THE TRANSVERSE DYNAMICS AND POLARIZATION OF GLUONS INSIDE THE PROTON AT THE LHC

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Transverse momentum dependent gluon distributions encode fundamental information on the structure of the proton. Here we show how they can be accessed in heavy quarkonium production in proton-proton collisions at the LHC. In particular, their first determination could come from the study of an isolated J/ψ or Υ particle, produced back to back with a photon.

1 Formalism

Transverse momentum dependent (TMD) gluon distributions inside an unpolarized proton are defined by the hadron matrix element of a correlator of the gluon field strengths $F^{\mu\rho}(0)$ and $F^{\nu\sigma}(\xi)$. Expanding the gluon four-momentum as $p = xP + p_T + p^-n$, with n being a lightlike vector conjugate to the momentum of the proton P , such correlator can be written as¹

$$\begin{aligned}\Phi_g^{\mu\nu}(x, \mathbf{p}_T) &= \frac{n_\rho n_\sigma}{(p \cdot n)^2} \int \frac{d(\xi \cdot P) d^2 \xi_T}{(2\pi)^3} e^{ip \cdot \xi} \langle P | \text{Tr} [F^{\mu\rho}(0) F^{\nu\sigma}(\xi)] | P \rangle \Big|_{\xi \cdot n=0} \\ &= -\frac{1}{2x} \left\{ g_T^{\mu\nu} f_1^g(x, \mathbf{p}_T^2) - \left(\frac{p_T^\mu p_T^\nu}{M_p^2} + g_T^{\mu\nu} \frac{\mathbf{p}_T^2}{2M_p^2} \right) h_1^{\perp g}(x, \mathbf{p}_T^2) \right\},\end{aligned}\quad (1)$$

where gauge links have been omitted. The transverse projector $g_T^{\mu\nu}$ is defined as $g_T^{\mu\nu} = g^{\mu\nu} - P^\mu n^\nu / P \cdot n - n^\mu P^\nu / P \cdot n$. Moreover, $p_T^2 = -\mathbf{p}_T^2$ and M_p is the proton mass. The gluon correlator of an unpolarized proton can therefore be expressed in terms of two independent TMD distribution functions: $f_1^g(x, \mathbf{p}_T^2)$ is the unpolarized one, while $h_1^{\perp g}(x, \mathbf{p}_T^2)$ denotes the T -even, helicity-flip distribution of linearly polarized gluons, which satisfies the model-independent positivity bound¹

$$\frac{\mathbf{p}_T^2}{2M_p^2} |h_1^{\perp g}(x, \mathbf{p}_T^2)| \leq f_1^g(x, \mathbf{p}_T^2).\quad (2)$$

Like any TMD distribution, $h_1^{\perp g}$ might receive contributions from initial and final state interactions that can render it nonuniversal and even hamper its extraction in processes for which TMD factorization does not apply.

2 Phenomenology

Several processes have been suggested to measure the experimentally unknown distributions f_1^g and $h_1^{\perp g}$. Although it has been discussed how to isolate the contribution from $h_1^{\perp g}$ by means of an azimuthal angular dependent weighting of the cross section for dijet production in hadronic collisions², TMD factorization is expected to be broken in this case due to the presence of both

initial and final state interactions³. A theoretically cleaner and safer way would be to study dijet or heavy quark pair production in electron-proton collisions, for instance at a future Electron-Ion Collider^{4,5}. Another process where the problem of factorization breaking is absent is $pp \rightarrow \gamma\gamma X$ ⁶, which however suffers from a huge background from π^0 decays and contaminations from quark-induced channels.

In the following we show how TMD gluon distributions can be probed in heavy quarkonium production at the LHC. TMD factorization should hold in this case, provided that the two quarks that form the bound state are produced in a colorless state already at short distances.

2.1 Transverse momentum distributions of $C = +$ quarkonia

We consider first the process $p(P_A) + p(P_B) \rightarrow \mathcal{Q}(\boldsymbol{\eta}) + X$, where \mathcal{Q} is a heavy quark-antiquark bound state with $C = +$, and the four-momenta of the particles are given between brackets. Assuming TMD factorization, the corresponding cross section can be written as

$$d\sigma = \frac{1}{2s} \frac{d^3\mathbf{q}}{(2\pi)^3 2\mathbf{q}^0} \int \mathbf{d}x_a dx_b d^2\mathbf{p}_{aT} d^2\mathbf{p}_{bT} (2\pi)^4 \delta^4(p_a + p_b - \boldsymbol{\eta}) \times \Phi_g^{\mu\nu}(x_a, \mathbf{p}_{aT}) \Phi_{\boldsymbol{\eta}}^{\rho\sigma}(x_b, \mathbf{p}_{bT}) \sum_{\text{colors}} \mathcal{A}_{\mu\rho} \mathcal{A}_{\nu\sigma}^*(p_a, p_b; \boldsymbol{\eta}), \quad (3)$$

with $s = (P_A + P_B)^2$ being the total energy squared in the hadronic center-of-mass frame and \mathcal{A} denoting the hard scattering amplitude of the dominant subprocess $g(p_a) + g(p_b) \rightarrow \mathcal{Q}(\boldsymbol{\eta})$. The amplitude \mathcal{A} is evaluated at order α_s^2 within the framework of the color-singlet model. Color octet contributions should be negligible, according to nonrelativistic QCD arguments⁷. For small transverse momentum, $\mathbf{q}_T^2 \ll M_{\mathcal{Q}}^2$, with $M_{\mathcal{Q}}$ being the quarkonium mass, the resulting transverse momentum distributions for η_Q and $\chi_{Q0,2}$ ($Q = c, b$) are given by

$$\begin{aligned} \frac{1}{\sigma(\eta_Q)} \frac{d\sigma(\eta_Q)}{dy d\mathbf{q}_T^2} &= \frac{\mathcal{C}[f_1^g f_1^g]}{\int d\mathbf{q}_T^2 \mathcal{C}[f_1^g f_1^g]} [1 - R(\mathbf{q}_T^2)], \\ \frac{1}{\sigma(\chi_{Q0})} \frac{d\sigma(\chi_{Q0})}{d\mathbf{y} d\mathbf{q}_T^2} &= \frac{\mathcal{C}[f_1^g f_1^g]}{\int d\mathbf{q}_T^2 \mathcal{C}[f_1^g f_1^g]} [1 + R(\mathbf{q}_T^2)], \\ \frac{1}{\sigma(\chi_{Q2})} \frac{d\sigma(\chi_{Q2})}{d\mathbf{y} d\mathbf{q}_T^2} &= \frac{\mathcal{C}[f_1^g f_1^g]}{\int d\mathbf{q}_T^2 \mathcal{C}[f_1^g f_1^g]}, \end{aligned} \quad (4)$$

where $\sigma = \int d\mathbf{q}_T^2 d\sigma$ and y is the rapidity of the quarkonium along the direction of the incoming protons. Furthermore, $x_{a,b} = M_{\mathcal{Q}}/\sqrt{s} e^{\pm y}$,

$$R(\mathbf{q}_T^2) = \frac{\mathcal{C}[w h_1^{\perp g} h_1^{\perp g}]}{\mathcal{C}[f_1^g f_1^g]}, \quad w = \frac{1}{2M^4} \left[(\mathbf{p}_{aT} \cdot \mathbf{p}_{bT})^2 - \frac{1}{2} \mathbf{p}_{aT}^2 \mathbf{p}_{bT}^2 \right], \quad (5)$$

and we have used the following definition of convolution of two TMD distributions f and g ,

$$\mathcal{C}[w f g] \equiv \int \mathbf{d}^2\mathbf{p}_{aT} \int \mathbf{d}^2\mathbf{p}_{bT} \delta^2(\mathbf{p}_{aT} + \mathbf{p}_{bT} - \mathbf{q}_T) w(\mathbf{p}_{aT}, \mathbf{p}_{bT}) f(x_a, \mathbf{p}_{aT}^2) g(x_b, \mathbf{p}_{bT}^2). \quad (6)$$

Our numerical estimates are shown in Fig.1, where we have assumed that the gluon distributions have a simple Gaussian dependence on transverse momentum. Namely,

$$f_1^g(x, \mathbf{p}_T^2) = \frac{f_1^g(x)}{\pi \langle p_T^2 \rangle} \exp\left(-\frac{\mathbf{p}_T^2}{\langle p_T^2 \rangle}\right), \quad (7)$$

where $f_1^g(x)$ is the collinear gluon distribution and the width $\langle p_T^2 \rangle$ is taken to be independent of x and the energy scale, set by $M_{\mathcal{Q}}$. The bound in Eq. (2) is satisfied, although not everywhere saturated, by the form

$$h_1^{\perp g}(x, \mathbf{p}_T^2) = \frac{M_p^2 f_1^g(x)}{\pi \langle p_T^2 \rangle^2} \frac{2(1-r)}{r} \exp\left(1 - \frac{1}{r} \frac{\mathbf{p}_T^2}{\langle p_T^2 \rangle}\right), \quad (8)$$

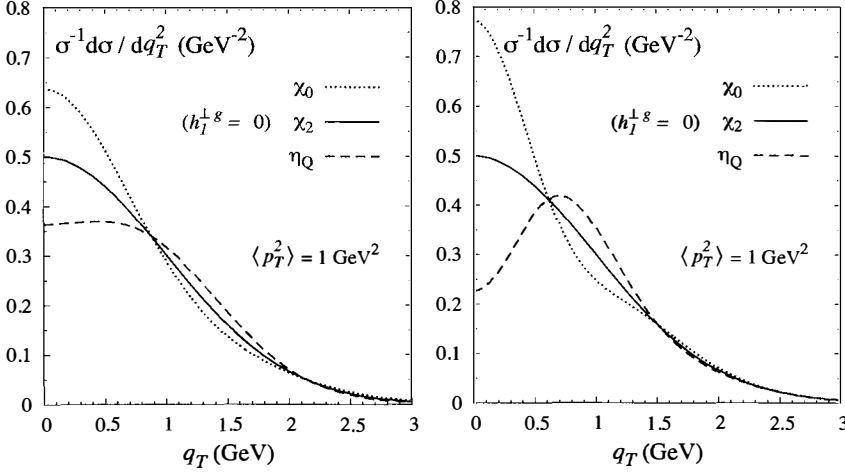


Figure 1 – Transverse momentum distributions for $C = +$ quarkonia evaluated at $y = 0$, obtained using the input distributions in Eqs. (7)–(8), with $\langle p_T^2 \rangle = 1 \text{ GeV}^2$ and two different values of r : $r = 2/3$ (left) and $r = 1/3$ (right).

with $0 < r < 1$. The distributions for $\eta_{c,b}$ and $\chi_{c,b0}$ are similar to the ones for a pseudoscalar and a scalar Higgs boson^{8,9}, and can be used to extract $h_1^{\perp g}$, while f_1^g can be accessed by looking at $\chi_{c,b2}$. A comparison among the different spectra could help to cancel out uncertainties. This experiment requires forward detectors like the LHCb, which hopefully will be able to provide such data in the near future.

2.2 $C = -$ quarkonium production in association with a photon

Along the lines of the previous section, we study the process $p(P_A) + p(P_B) \rightarrow Q(P_Q) + \gamma(P_\gamma) + X$, where now Q is a $C = -$ quarkonium (J/ψ or Υ) produced almost back to back with the photon. Hence the imbalance $\mathbf{q}_T = \mathbf{P}_{Q_T} + \mathbf{P}_{\gamma_T}$ will be small, but not the individual transverse momenta of the two particles. No forward detector is therefore needed in this case. The cross section has the following structure,

$$\frac{d\sigma}{dQ dY d^2\mathbf{q}_T d\Omega} \propto F_1 \mathcal{C} [f_1^g f_1^g] + F_3 \mathcal{C} [w_3 f_1^g h_1^{\perp g} + x_a \leftrightarrow x_b] \cos 2\phi + F_4 \mathcal{C} [w_4 h_1^{\perp g} h_1^{\perp g}] \cos 4\phi, \quad (9)$$

where Q and Y are the invariant mass and the rapidity of the pair, to be measured, like \mathbf{q}_T , in the hadronic center-of-mass frame. On the other hand, the solid angle $\Omega = (\theta, \phi)$ is measured in the Collins-Soper frame, where the final pair is at rest and the $\hat{x}\hat{z}$ -plane is spanned by \mathbf{P}_A and \mathbf{P}_B , with the \hat{x} -axis set by their bisector. The transverse weights are given by

$$w_3 = \frac{\mathbf{q}_T^2 \mathbf{P}_{bT}^2 - 2(\mathbf{q}_T \cdot \mathbf{P}_{bT})^2}{2M_p^2 \mathbf{q}_T^2}, \quad w_4 = 2 \left[\frac{\mathbf{P}_{aT} \cdot \mathbf{P}_{bT}}{2M_p^2} - \frac{(\mathbf{P}_{aT} \cdot \mathbf{q}_T)(\mathbf{P}_{bT} \cdot \mathbf{q}_T)}{M_p^2 \mathbf{q}_T^2} \right]^2 - \frac{\mathbf{P}_{aT}^2 \mathbf{P}_{bT}^2}{4M_p^4}, \quad (10)$$

and the light-cone momentum fractions are $x_{a,b} = \exp[\pm Y] Q/\sqrt{s}$. Explicit expressions for $F_{1,3,4}$ can be found elsewhere¹⁰. We propose the measurement of the following three observables,

$$S_{q_T}^{(n)} \equiv \frac{\int d\phi \cos(n\phi) \frac{d\sigma}{dQ dY d^2\mathbf{q}_T d\Omega}}{\int d\mathbf{q}_T^2 \int d\phi \frac{d\sigma}{dQ dY d^2\mathbf{q}_T d\Omega}}, \quad (11)$$

with $n = 0, 2, 4$, and where the \mathbf{q}_T^2 integration in the denominator is up to $(Q/2)^2$. In this way we are able to single out the three terms in Eq. 9:

$$\mathcal{S}_{q_T}^{(0)} = \frac{C[f_1^g f_1^g]}{\int d\mathbf{q}_T^2 C[f_1^g f_1^g]}, \quad \mathcal{S}_{q_T}^{(2)} = \frac{F_3 C[w_3 f_1^g h_1^{\perp g} + x_a \leftrightarrow x_b]}{2F_1 \int d\mathbf{q}_T^2 C[f_1^g f_1^g]}, \quad \mathcal{S}_{q_T}^{(4)} = \frac{F_4 C[w_4 h_1^{\perp g} h_1^{\perp g}]}{2F_1 \int d\mathbf{q}_T^2 C[f_1^g f_1^g]}. \quad (12)$$

Our model predictions are presented in Fig. 2 for $\Upsilon + \gamma$ production, in a kinematic region where color octet contributions are suppressed¹⁰. The size of $\mathcal{S}_{q_T}^{(0)}$ should be sufficient to allow for a determination of the shape of f_1^g as a function of \mathbf{q}_T . Since $\mathcal{S}_{q_T}^{(2)}$ and $\mathcal{S}_{q_T}^{(4)}$ are considerably smaller, one would need to integrate them over \mathbf{q}_T^2 [up to $(Q/2)^2$] to get at least an experimental evidence of a nonzero $h_1^{\perp g}$.

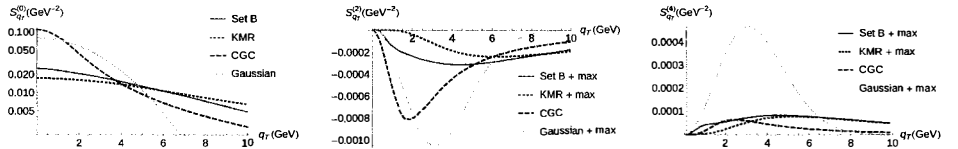


Figure 2 – Model predictions for $\mathcal{S}_{q_T}^{(0)}$, $\mathcal{S}_{q_T}^{(2)}$ and $\mathcal{S}_{q_T}^{(4)}$ for the process $p(P_A) + p(P_B) \rightarrow Q(P_Q) + \gamma(P_\gamma) + X$ at $\sqrt{s} = 14$ TeV in the kinematic region defined by $Q = 20$ GeV, $Y = 0$, $\theta = \pi/2$, and $x_a = x_b \simeq 1.4 \times 10^{-3}$.

3 Conclusions

The distribution of linearly polarized gluons inside an unpolarized proton $h_1^{\perp g}$ leads to a modulation of the transverse momentum distribution of scalar (χ_{c0} , χ_{b0}) and pseudoscalar (η_c , η_b) quarkonia that depends on their parity. It does not contribute to the transverse spectra of χ_{c2} and χ_{b2} , which can be used to probe the unpolarized gluon distribution f_1^g . No angular analysis is needed for such measurements and experimental opportunities are offered by LHCb and the proposed fixed-target experiment AFTER at LHC¹¹. Furthermore, a first determination of $h_1^{\perp g}$ and f_1^g could come from $J/\psi(\Upsilon) + \gamma$ production at the running experiments at the LHC, where yields are large enough to perform these analyses with existing data at $\sqrt{s} = 7$ and 8 TeV. We have shown that, together with similar studies in the Higgs sector, quarkonium production can be used to extract gluon TMDs and investigate their process and energy scale dependences.

Acknowledgments

This work was supported by the European Community under the Ideas program QWORK (contract 320389).

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