PROCEEDINGS OF SCIENCE

Studying the gluon TMDs with J/ψ- and Υ-pair production at the LHC

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We report on how J/ψ- and Υ-pair production are promising processes to access the polarised and unpolarised gluon TMDs at the LHC. We present the formalism used, as well as resulting observables that could be extracted from data.

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1. Introduction

Transverse-Momentum Dependent (TMD) factorisation allows one to study the correlations between the parton spin and its transverse momentum, inside polarised and unpolarised hadrons [1, 2, 3]. These correlations can manifest themselves through azimuthal modulations of the cross-sections for various hadronic processes. At the LHC, the gluon density inside the proton is much higher than the quark ones, making gluon fusion the main channel of heavy-flavoured hadron production. Quarkonium-pair production is a particularly interesting process for the study of gluon TMDs. The quarkonia are dominantely produced via colour-singlet transitions [4, 5, 6, 7]. The observed final state is thus colourless which is a necessary condition not to break TMD factorisation [8, 9]. Moreover, a two-particle final state allows for large individual momenta for each quarkonium adding up to a small transverse momentum of the pair, a requirement for the use of TMD factorisation.

J/ψ- and Υ- pair production has already been the object of several experimental studies at the LHC and the Tevatron [10, 11, 12, 13, 14, 15]. This allowed to perform a first fit of the width of a Gaussian-modelled TMD using the normalised transverse-momentum-spectrum of di-J/ψ production at LHCb hinting at the presence of QCD evolution effects [16]. More observables could be extracted from data already acquired and to come and this motivates us advancing the theory description of such processes by including TMD QCD evolution as we report on here along the lines of [17].

2. Quarkonium-pair production within TMD factorisation

TMD factorisation extends collinear factorisation by accounting for the parton transverse momenta. The hard-scattering amplitude is factorised into a short-distance coefficient and parton correlators that contain the dependence on the partonic transverse momentum. The gluon correlator for an unpolarised proton can be parametrised in terms of two independent TMDs [18, 19]. The first distribution is that of unpolarised gluons $f_1^g(x,k_T^2)$, the second one is that of linearly polarised gluons $h_1^{\perp g}(x,k_T^2)$. The generic structure of the TMD cross-section for quarkonium-pair production from gluon fusion is similar to that of quarkonium-photon production [20, 21] and reads [16]:

$$\frac{\mathrm{d}\sigma}{\mathrm{d}M_{QQ}\mathrm{d}Y_{QQ}\mathrm{d}^2\boldsymbol{P}_{QQ_T}\mathrm{d}\Omega} = \frac{\sqrt{M_{QQ}^2-4M_Q^2}}{(2\pi)^2 8s\,M_{QQ}^2}\Big\{F_1\, \mathcal{C}\Big[f_1^g f_1^g\Big] + F_2\, \mathcal{C}\Big[w_2 h_1^{\perp g} h_1^{\perp g}\Big]$$
$$+\cos 2\phi_{\mathrm{CS}}\left(F_3\, \mathcal{C}\Big[w_3 f_1^g h_1^{\perp g}\Big] + F_3'\, \mathcal{C}\Big[w_3' h_1^{\perp g} f_1^g\Big]\right) + \cos 4\phi_{\mathrm{CS}}\, F_4\, \mathcal{C}\Big[w_4 h_1^{\perp g} h_1^{\perp g}\Big]\Big\}, \tag{2.1}$$

with $\mathcal{C}[w\,f\,g](x_{1,2},\boldsymbol{P}_{QQ_T}) \equiv \int \mathrm{d}^2k_{1T}\int \mathrm{d}^2k_{2T}\,\delta^2(\boldsymbol{k}_{1T}+\boldsymbol{k}_{2T}-\boldsymbol{P}_{QQ_T})\,w(\boldsymbol{k}_{1T},\boldsymbol{k}_{2T})\,f(x_1,\boldsymbol{k}_{1T}^2)\,g(x_2,\boldsymbol{k}_{2T}^2)$ (the w_i are called TMD weights), $\mathrm{d}\Omega = \mathrm{d}\cos\theta_{\mathrm{CS}}\mathrm{d}\phi_{\mathrm{CS}}$, $\{\theta_{\mathrm{CS}},\phi_{\mathrm{CS}}\}$ being the Collins-Soper (CS) angles [22], Y_{QQ} is the pair rapidity and $s=(P_1+P_2)^2$. $\boldsymbol{P}_{QQ_T}$ and Y_{QQ} are considered in the hadron c.m.s. The F_i coefficients are functions of θ_{CS} and the invariant mass of the pair, denoted M_{QQ}.

The weights in Eq. (2.1) are identical for all gluon-fusion processes in unpolarised proton collisions and can be found in [21]. Having at hand the hard-scattering coefficients, one can extract the TMD convolutions from measurements of the cross-section azimuthal modulations. To do so, one defines $\cos(n\phi_{\mathrm{CS}})$-weighted differential cross-sections, integrated over ϕ_{CS} and normalized by their azimuthally-independent component:

$$\langle\cos(n\phi_{\rm CS})\rangle = \frac{\displaystyle\int d\phi_{\rm CS}\cos(n\phi_{\rm CS})\,\frac{d\sigma}{dM_{QQ}dY_{QQ}d^2\boldsymbol{P}_{QQ_T}d\Omega}}{\displaystyle\int d\phi_{\rm CS}\,\frac{d\sigma}{dM_{QQ}dY_{QQ}d^2\boldsymbol{P}_{QQ_T}d\Omega}}\,. \tag{2.2}$$

The TMD QCD evolution introduces a dependence of the TMDs on two scales: a renormalization scale μ and a rapidity scale ζ [23, 24, 25, 26]. The evolution is most easily treated in the impact-parameter space, b_T being the conjugate variable to k_T. TMDs are first computed at the scale $\mu \sim \sqrt{\zeta} \sim \mu_b = b_0/b_T$ (with $b_0 = 2e^{-\gamma_E}$) to minimise the logarithms of μb_T and ζb_T^2, and then evolved up to $\mu \sim \sqrt{\zeta} \sim M_{QQ}$, the natural scale of the hard part. Solving the renormalisation group equation and the Collins-Soper equation one can show that such evolution introduces a Sudakov factor S_A in the convolution. This factor can be perturbatively evaluated for small values of b_T, and is given at next-to-leading-logarithmic accuracy by [26]:

$$S_A(b_T;\zeta,\mu) = \frac{C_A}{\pi}\int_{\mu_b^2}^{\mu^2}\frac{d\bar{\mu}^2}{\bar{\mu}^2}\alpha_s(\bar{\mu}^2)\left[\left(1+\frac{\alpha_s(\bar{\mu}^2)}{4\pi}\,\frac{67-3\pi^2-20T_f n_f}{9}\right)\ln\left(\frac{\zeta}{\bar{\mu}^2}\right)-\frac{11-2n_f/C_A}{6}\right]. \tag{2.3}$$

In order to perform the Fourier transform, one needs to integrate over all b_T values inside the convolution while paying attention to the validity range of Eq. (2.3). Indeed, large b_T values correspond to the non-perturbative region, and too small b_T values would correspond to $\mu_b > M_{QQ}$ for which the expression is not valid anymore. The TMD scale $\mu_b(b_T)$ is thus replaced by $\mu_b\big(b_T^*(b_c(b_T))\big)$ following [27] such that μ_b is bound to lie between $b_0/b_{T_{\rm max}}$ and M_{QQ}.

As far as the TMD expressions are concerned, they can be pertubatively computed at leading order in α_s as

$$\tilde{f}_1^g(x, b_T^{*2};\zeta,\mu) = f_{g/P}(x;\mu) + O(\alpha_s), \tag{2.4}$$

$$\tilde{h}_1^{\perp g}(x, b_T^{*2};\zeta,\mu) = \frac{\alpha_s(\mu)}{\pi}\int_x^1\frac{d\hat{x}}{\hat{x}}\left(\frac{\hat{x}}{x}-1\right)\left(C_A f_{g/P}(\hat{x};\mu) + C_F\sum_{i=q,\bar{q}} f_{i/P}(\hat{x};\mu)\right) + O(\alpha_s^2)\,. \tag{2.5}$$

As $h_1^{\perp g}$ describes the correlation between the gluon polarisation and its transverse momentum k_T inside the unpolarised proton, it requires a helicity flip and thus an additional gluon emission. Consequently, its perturbative expansion starts at $O(\alpha_s)$ [28]. Such a reasoning only applies for b_T in the perturbative region, outside of which $h_1^{\perp g}$ is not computable.

One still needs to describe the large-b_T behaviour of both the initial TMDs and the perturbative Sudakov factor. The deviation between the components perturbatively evaluated at b_T^* and their actual value is encoded in what is called a nonperturbative Sudakov factor $S_{\rm NP}$. Including all these ingredients, one gets the expressions for the different TMD convolutions in b_T-space to be:

$$\mathcal{C}\Big[f_1^g f_1^g\Big] = \int_0^\infty\frac{db_T^2}{4\pi}\,J_0(b_T q_T)\,e^{-S_A(b_T^*;M_{QQ}^2,M_{QQ})-S_{\rm NP}(b_c)}\,\tilde{f}_1^g(x_1,b_T^{*2};\mu_b^2,\mu_b)\,\tilde{f}_1^g(x_2,b_T^{*2};\mu_b^2,\mu_b) \tag{2.6}$$

$$\mathcal{C}\Big[w_2\,h_1^{\perp g}h_1^{\perp g}\Big] = \int_0^\infty\frac{db_T^2}{4\pi}\,J_0(b_T q_T)\,e^{-S_A(b_T^*;M_{QQ}^2,M_{QQ})-S_{\rm NP}(b_c)}\,\tilde{h}_1^{\perp g}(x_1,b_T^{*2};\mu_b^2,\mu_b)\,\tilde{h}_1^{\perp g}(x_2,b_T^{*2};\mu_b^2,\mu_b)\,,$$

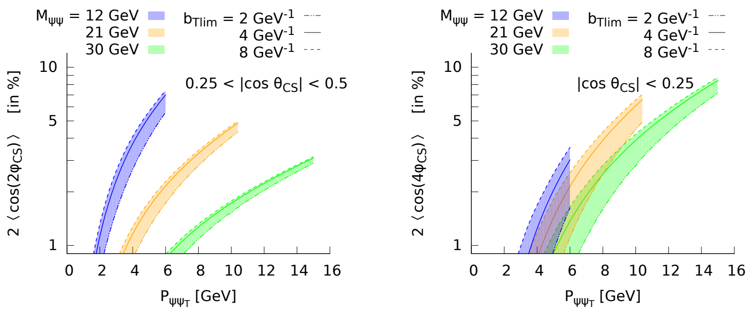
$$\mathcal{C}\Big[w_3\,f_1^g h_1^{\perp g}\Big] = \int_0^\infty\frac{db_T^2}{4\pi}\,J_2(b_T q_T)\,e^{-S_A(b_T^*;M_{QQ}^2,M_{QQ})-S_{\rm NP}(b_c)}\,\tilde{f}_1^g(x_1,b_T^{*2};\mu_b^2,\mu_b)\,\tilde{h}_1^{\perp g}(x_2,b_T^{*2};\mu_b^2,\mu_b)\,,$$

$$\mathcal{C}\Big[w_4\,h_1^{\perp g}h_1^{\perp g}\Big] = \int_0^\infty\frac{db_T^2}{4\pi}\,J_4(b_T q_T)\,e^{-S_A(b_T^*;M_{QQ}^2,M_{QQ})-S_{\rm NP}(b_c)}\,\tilde{h}_1^{\perp g}(x_1,b_T^{*2};\mu_b^2,\mu_b)\,\tilde{h}_1^{\perp g}(x_2,b_T^{*2};\mu_b^2,\mu_b)\,.$$

This $S_{\rm NP}$ factor is unknown. We have used [17] a very simple form for it in order to estimate its impact over TMD observables, namely $S_{\rm NP}(b_c(b_T)) = A \ln\left(\frac{M_{QQ}}{Q_{\rm NP}}\right) b_c^2(b_T)$ with $Q_{\rm NP} = 1$ GeV where $Q_{\rm NP}$ is a reference scale and A the only free parameter in our model that parametrises the width of the b_T-Gaussian. We have performed computations for values of the parameter A equal to 0.64, 0.16 and 0.04 GeV^2, which respectively correspond to a factor $e^{-S_{\rm NP}}$ becoming smaller than 10^{-3} at $b_{T\,\rm lim} = 2$, 4 and 8 GeV^{-1}.

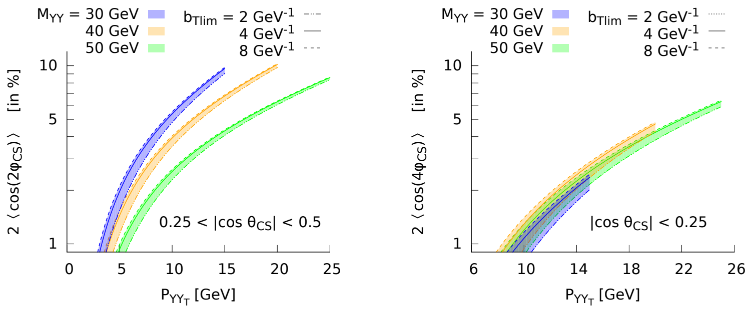
3. Results for J/ψ- and Υ-pair production

Using Eq. (2.6) and knowing the F_i, one can compute the $\cos(2,4\phi)$-asymmetries. In Fig. 1, we present the asymmetries for di-J/ψ production as functions of $\boldsymbol{P}_{QQ_T}$. One can see that the asymmetries grow rapidly with $\boldsymbol{P}_{QQ_T}$, as the convolutions containing $h_1^{\perp g}$ have a harder $\boldsymbol{P}_{QQ_T}$-spectrum than $\mathcal{C}\Big[f_1^g f_1^g\Big]$. The $M_{\psi\psi}$-dependence is mostly coming from the hard-scattering coefficients. We observe that the uncertainty band generated by the variation of the width of $S_{\rm NP}$ is narrower at large scales, where the nonperturbative component of the TMDs is less relevant. The size of the asymmetries can reach up to 10%. Fig. 2 displays the same asymmetries for di-Υ production, at typically higher energies. At these scales, the $S_{\rm NP}$-width uncertainty is quite narrow.



(a) (b)

Figure 1: The azimuthal asymmetries for di-J/ψ production as functions of $\boldsymbol{P}_{QQ_T}$. The different plots show $2\langle\cos(2\phi_{CS})\rangle$ at $0.25 < |\cos(\theta_{CS})| < 0.5$ (a) and $2\langle\cos(4\phi_{CS})\rangle$ at $|\cos(\theta_{CS})| < 0.25$ (b). Results are presented for $M_{\psi\psi} = 12$, 21 and 30 GeV, and for $b_{T\,\rm lim} = 2$, 4 and 8 GeV^{-1}.



(a) (b)

Figure 2: The azimuthal asymmetries for di-Υ production as functions of $\boldsymbol{P}_{QQ_T}$. The different plots show $2\langle\cos(2\phi_{CS})\rangle$ at $0.25 < |\cos(\theta_{CS})| < 0.5$ (a) and $2\langle\cos(4\phi_{CS})\rangle$ at $|\cos(\theta_{CS})| < 0.25$ (b). Results are presented for $M_{\Upsilon\Upsilon} = 30$, 40 and 50 GeV, and for $b_{T\,\rm lim} = 2$, 4 and 8 GeV^{-1}. Results for $M_{\Upsilon\Upsilon} = 30$ GeV are not included in Fig. 2d as they are below percent level.

4. Conclusions

Quarkonium-pair production offers great opportunities to extract the gluon TMDs inside unpolarised protons. We reported on the formalism used to estimate the size of the related asymmetries, including TMD evolution in the computations. The full analysis can be found in [17]. We note that more efforts are needed to obtain rigorous factorisation theorems along the lines of [29]. Several asymmetries could be observed at the LHC, both in di-J/ψ and di-Υ production. Accessing these poorly known distributions would improve our understanding of the proton structure, in addition to other TMDs.

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