IL NUOVO CIMENTO **49 C** (2026) 41
DOI 10.1393/ncc/i2026-26041-4

Colloquia: IFAE 2025

Gluon TMDs from C-even quarkonium production(*)

Nanako Kato(1)(2)

(1) *INFN, Sezione di Cagliari - Cagliari, Italy*
(2) *Dipartimento di Fisica, Università degli Studi di Cagliari - Cagliari, Italy*

received 10 October 2025

Summary. — We present the numerical results for the upper bounds of the single-spin asymmetries for the production of quarkonium states with even charge conjugation in proton-proton collisions. The theoretical framework adopted for the description of quarkonium production is non-relativistic QCD. The observables computed in this study could be measured at LHCSpin, the fixed-target experiment planned at the LHC.

1. – Introduction

One of the most promising processes that can probe the gluon content of the proton is the production of bound states of heavy quarks (quarkonia) in proton-proton collisions. In particular, these processes can provide experimental observables sensitive to gluon transverse momentum dependent parton distribution functions (TMD PDFs), that encode information on the intrinsic motion of gluons and the correlations between their spins and momenta, providing a full three-dimensional picture of hadrons.

In this contribution, we focus on the production of scalar and pseudoscalar quarkonia with even charge conjugation (C-even), characterized by total angular momentum J, parity P and charge conjugation $J^{PC} = 0^{\pm +}$. Using the spectroscopic notation $^{2S+1}L_J$, with S being the spin and L the orbital angular momentum, we study the following quarkonium states: $^1S_0$ states η_c, η_b, the $^3P_0$ states χ_{c0}, χ_{b0}, and the $^3P_2$ states χ_{c2}, χ_{b2}. Thanks to the charm and bottom masses being large, perturbative QCD remains applicable even if the transverse momentum q_T of the quarkonium state is small ($q_T^2 \ll M_{\mathcal{Q}}^2$). In this kinematic regime, TMD factorization is expected to be applicable. Moreover, these quarkonium states are produced, at leading order in the strong coupling constant α_s, by gluon-gluon fusion without additional gluon emission in the final state.

We adopt the effective field theory approach of nonrelativistic QCD (NRQCD) [1], according to which quarkonium production in proton-proton collisions is described by

(*) IFAE 2025 - "Energy Frontier" session

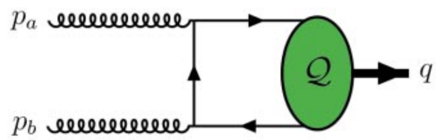

Fig. 1. – Leading order diagram for the process $gg \to \mathcal{Q}$, where $\mathcal{Q}$ is the heavy quark-antiquark bound state. The crossed diagram, in which the directions of the arrows in the fermionic lines are reversed, is not shown.

a double power series expansion in α_s and the relative velocity v of the heavy quark-antiquark pair in the quarkonium rest frame, with $v \ll 1$. The hadronization of the pairs, which can be produced both in color-singlet and color-octet states, is encoded in the long distance matrix elements (LDMEs). These are non-perturbative objects we need to extract from experimental data. Since for the quarkonium states under study the color-octet contributions are suppressed, we will assume that the heavy quark-antiquark pairs are produced directly with the same quantum numbers as the observed bound states, as in the traditional color-singlet model (CSM) [2].

The observables investigated here could, in principle, be measured by the LHCSpin project, a proposed polarized gas target experiment at the LHCb spectrometer [3, 4].

2. – Cross sections and azimuthal modulations

We consider the inclusive quarkonium production in a proton-proton collision process [5, 6]

$$p(P_A, S_A) + p(P_B, S_B) \to Q\bar{Q}[^{2S+1}L_J^{(1)}](q) + X, \tag{1}$$

where P_i and S_i, with $i = A, B$, represent, respectively, the four-momenta and the spin of the colliding protons, q is the four-momentum of the heavy quark-antiquark pair ($Q\bar{Q}$), produced in an intermediate Fock state $^{2S+1}L_J^{(1)}$, in a colorless configuration specified by the superscript (1). The squared invariant mass of the resonance is $M^2 = q^2$, with M being twice the heavy quark mass.

At the lowest order in perturbative QCD, the partonic reaction to be considered is the $2 \to 1$ gluon-gluon fusion process

$$g(p_a) + g(p_b) \to Q\bar{Q}[^{2S+1}L_J^{(1)}](q), \tag{2}$$

as depicted in fig. 1.

The gluon momenta can be approximated as follows

$$p_a = x_a P_A + p_{aT}\,, \qquad p_b = x_b P_B + p_{bT}\,, \tag{3}$$

in terms of the longitudinal momentum fractions (x_a, x_b) and the transverse momenta (p_{aT}, p_{bT}). In the kinematic region where the transverse momentum q_T of the produced gluon pol.

nucleon pol.	U	circ.	lin.
U	f_1^g		$h_1^{\perp g}$
L		g_1^g	$h_{1L}^{\perp g}$
T	$f_{1T}^{\perp g}$	g_{1T}^g	$h_1^g,\, h_{1T}^{\perp g}$

Fig. 2. – Twist-2 gluon TMDs. U, L, T correspond to unpolarized, longitudinally polarized and transversely polarized nucleons. U, circ., lin. correspond to unpolarized, circularly polarized and linearly polarized gluons. Functions in blue are T-even.

quarkonium state is much smaller than its invariant mass, namely $q_T \ll M$, the cross section for the process in eq. (1) is found to be

(4)
$$\frac{d\sigma[\mathcal{Q}]}{dy\, d^2\boldsymbol{q}_T} = F_{UU}^{\mathcal{Q}} + F_{UL}^{\mathcal{Q}}\, S_{BL} + F_{LU}^{\mathcal{Q}}\, S_{AL} + F_{UT}^{\mathcal{Q},\sin\phi_{S_B}}\, |\boldsymbol{S}_{BT}| \sin\phi_{S_B} \\ + F_{TU}^{\mathcal{Q},\sin\phi_{S_A}}\, |\boldsymbol{S}_{AT}| \sin\phi_{S_A} + F_{LL}^{\mathcal{Q}}\, S_{AL}\, S_{BL} + F_{LT}^{\mathcal{Q},\cos\phi_{S_B}}\, S_{AL}\, |\boldsymbol{S}_{BT}| \cos\phi_{S_B} \\ + F_{TL}^{\mathcal{Q},\cos\phi_{S_A}}\, |\boldsymbol{S}_{AT}|\, S_{BL} \cos\phi_{S_A} + |\boldsymbol{S}_{AT}|\, |\boldsymbol{S}_{BT}| \Big(F_{TT}^{\mathcal{Q},\cos(\phi_{S_A}-\phi_{S_B})} \cos(\phi_{S_A}-\phi_{S_B}) \\ + F_{TT}^{\mathcal{Q},\cos(\phi_{S_A}+\phi_{S_B})} \cos(\phi_{S_A}+\phi_{S_B}) \Big)\,,$$

with y being the rapidity of the outgoing quarkonium. Furthermore, ϕ_{S_A} (ϕ_{S_B}) is the azimuthal angle of the spin vector S_A (S_B). The subscripts of the structure functions $F^{\mathcal{Q}}$ refer to the polarization of the incoming protons. Each structure function in eq. (4) can be factorized into a hard part $H^{\mathcal{Q}}$, which is calculable as a perturbative expansion in α_s,

(5)
$$H^{\mathcal{Q}} = \sum_{n=0}^{\infty} \left(\frac{\alpha_s}{\pi}\right)^n H_n^{\mathcal{Q}}\,,$$

and a non-perturbative part, given by one or more convolutions of gluon TMDs multiplied by the proper LDME. In momentum space, these convolutions are defined as

(6)
$$\mathcal{C}[w\, f_1^g\, f_2^g] = \int d^2\boldsymbol{p}_{aT}\, d^2\boldsymbol{p}_{bT}\, w(\boldsymbol{p}_{aT}, \boldsymbol{p}_{bT})\, f_1^g(x_a, \boldsymbol{p}_{aT})\, f_2^g(x_b, \boldsymbol{p}_{bT})\, \delta^2(\boldsymbol{p}_{aT} + \boldsymbol{p}_{bT} - \boldsymbol{q}_T)\,,$$

where f_i, with $i = 1, 2$, are the gluon TMDs and $w(\boldsymbol{p}_{aT}, \boldsymbol{p}_{bT})$ is a proper weight function that depends on the particular gluon distributions involved. In this study, we take into account only the eight leading-twist gluon TMDs listed in fig. 2.

Considering the configuration where the proton with momentum P_A is unpolarized and the proton with momentum P_B is transversely polarized, we can define the azimuthal moments as

(7)
$$A_N^{\mathcal{Q},\sin\phi_{S_B}} = 2\, \frac{\int d\phi_{S_B} \sin\phi_{S_B}\, [d\sigma(\phi_{S_B}) - d\sigma(\phi_{S_B}+\pi)]}{\int d\phi_{S_B}\, [d\sigma(\phi_{S_B}) + d\sigma(\phi_{S_B}+\pi)]} = \frac{F_{UT}^{\mathcal{Q},\sin\phi_{S_B}}}{F_{UU}^{\mathcal{Q}}}$$

We find that the azimuthal moments can be written in terms of ratios of convolutions containing the gluon TMDs, as follows

$$
\begin{aligned}
A_N^{\eta_Q,\sin\phi_{S_B}} &= \frac{-\mathcal{C}[w_{UT}^f f_1^g f_{1T}^{\perp g}] + \mathcal{C}[w_{UT}^h h_1^{\perp g} h_1^g] - \mathcal{C}[w_{UT}^{h\perp} h_1^{\perp g} h_1^{\perp g}]}{\mathcal{C}[f_1^g f_1^g] - \mathcal{C}[w_{UU}^h h_1^{\perp g} h_1^{\perp g}]}, \\
A_N^{\chi_{Q0},\sin\phi_{S_B}} &= \frac{-\mathcal{C}[w_{UT}^f f_1^g f_{1T}^{\perp g}] - \mathcal{C}[w_{UT}^h h_1^{\perp g} h_1^g] + \mathcal{C}[w_{UT}^{h\perp} h_1^{\perp g} h_1^{\perp g}]}{\mathcal{C}[f_1^g f_1^g] + \mathcal{C}[w_{UU}^h h_1^{\perp g} h_1^{\perp g}]}, \\
A_N^{\chi_{Q2},\sin\phi_{S_B}} &= -\frac{\mathcal{C}[w_{UT}^f f_1^g f_{1T}^{\perp g}]}{\mathcal{C}[f_1^g f_1^g]},
\end{aligned}
\tag{8}
$$

where the explicit expressions of the weight functions can be found in [6].

Such a process could be in principle accessible at LHCSpin, the fixed target experiment planned at the LHC.

3. – Numerical results

In this section, we present the phenomenological estimates of the upper bounds of the single spin asymmetries (SSAs) in eq. (8), by assuming that the unpolarized gluon TMD has the following Gaussian form [7, 8]:

$$
f_1^g(x, \boldsymbol{p}_T^2) = \frac{f_1^g(x)}{\pi\langle p_T^2\rangle} \exp\left[-\frac{\boldsymbol{p}_T^2}{\langle p_T^2\rangle}\right], \tag{9}
$$

with $f_1^g(x)$ being the collinear gluon unpolarized distribution. The effect of the other unknown TMDs will be maximal when they saturate the following, model-independent, positivity bounds [5, 9]

$$
\begin{aligned}
\left|f_{1T}^{\perp g}(x, \boldsymbol{p}_T^2)\right|,\ \left|h_1^g(x, \boldsymbol{p}_T^2)\right| &\le \frac{M_p}{|\boldsymbol{p}_T|} f_1^g(x, \boldsymbol{p}_T^2), \\
\frac{1}{2}\left|h_1^{\perp g}(x, \boldsymbol{p}_T^2)\right| &\le \frac{M_p^2}{\boldsymbol{p}_T^2} f_1^g(x, \boldsymbol{p}_T^2), \\
\frac{1}{2}\left|h_{1T}^{\perp g}(x, \boldsymbol{p}_T^2)\right| &\le \frac{M_p^3}{|\boldsymbol{p}_T|^3} f_1^g(x, \boldsymbol{p}_T^2).
\end{aligned}
\tag{10}
$$

We show the results on the upper bounds of the azimuthal moments in eq. (8) for the $J = 0$ and $J = 2$ charmonium states in fig. 3 as a function of q_T. Existing phenomenological analyses show that the value $\langle p_T^2\rangle = 1$ GeV^2 should be a reasonable choice for the Gaussian width of the unpolarized gluon TMD at the scale $\mu^2 = 4M_c^2$ [10-12]. Moreover, we show our prediction for $q_T \le 2$ GeV, to ensure that the analysis remains within the kinematic region where TMD factorization is expected to be applicable. The parameters ρ_i ($0 < \rho_i < 1$), that enter the widths of the Gaussian parametrization of the polarized gluon TMDs [6], have been chosen in order to maximize the asymmetries.

In fig. 4 we show the upper bounds of the azimuthal moments in eq. (8), this time taking $\langle p_T^2\rangle = 3$ GeV^2, since larger values of $\langle p_T^2\rangle$ are expected for bottomonium production and the energy scale of this process would be similar to that explored in di-J/ψ production in ref. [8]. Moreover, since the transverse momentum region where the TMD

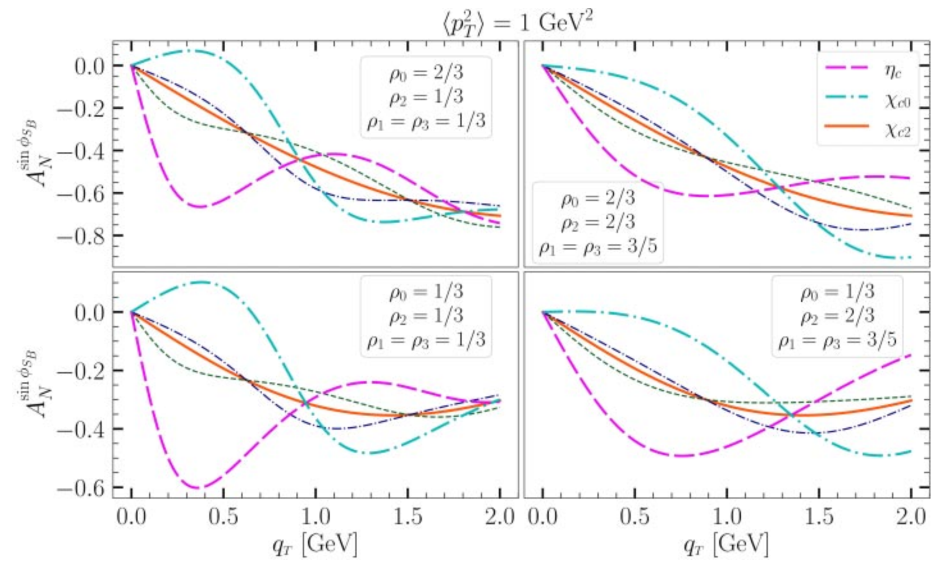

Fig. 3. – SSAs for different C-even quarkonia as a function of q_T. More specifically, the orange solid line corresponds to χ_{Q2}, the magenta long-dashed line to η_Q, and the light-blue dash-dotted line to χ_{Q0} states. The gluon TMDs are evaluated for different choices of the parameters ρ_i (check the text boxes in each panel for more details) entering the adopted Gaussian parametrization. The dark blue dashed and dark green dash-dotted thin lines represent the SSAs of η_Q and χ_{Q0}, respectively, with the TMDs h_1^g and $h_{1T}^{\perp g}$ set to 0.

factorization is applicable becomes wider for higher scales, we choose $q_T = 3$ GeV. From a direct comparison between figs. 3 and 4, we see that the predictions driven by the Gaussian parameterizations of the TMDs are qualitatively similar, exhibiting the same features but at different values of q_T. More specifically, the SSAs at $\langle p_T^2 \rangle = 3$ GeV^2 are broader in q_T, as expected from the nature of the Gaussian parameterizations employed. Analytical expressions for the SSAs show that for the production of χ_{Q2}, they are totally driven by the gluon Sivers function $f_{1T}^{\perp g}$, which describes the distributions of unpolarized gluons inside a transversely polarized nucleon. The SSAs for χ_{Q0} and η_Q, on the other hand, are sensitive to the linearly polarized gluon TMDs, h_1^g and $h_{1T}^{\perp g}$, and to the Sivers TMD as well. Therefore, by comparing the SSAs for these quarkonium states with those for χ_{Q2} one should be able to disentangle the combined effects of the linearly polarized gluon TMDs. According to the Gaussian model, the maximal impact of h_1^g and $h_{1T}^{\perp g}$ occurs for $q_T \leq 1$ GeV, suggesting that this kinematic region will be particularly important for accessing these completely unknown gluon TMDs.

4. – Conclusions

In this contribution, we have investigated the production of C-even quarkonium states in (un)polarized proton-proton collisions, demonstrating how these processes can be used to probe the internal structure of the nucleon through the extraction of gluon TMDs, which are still unknown from the experimental side. We have derived analytical expressions for the cross section within the framework of transverse momentum dependent factorization, in combination with NRQCD. Furthermore, we have provided numerical

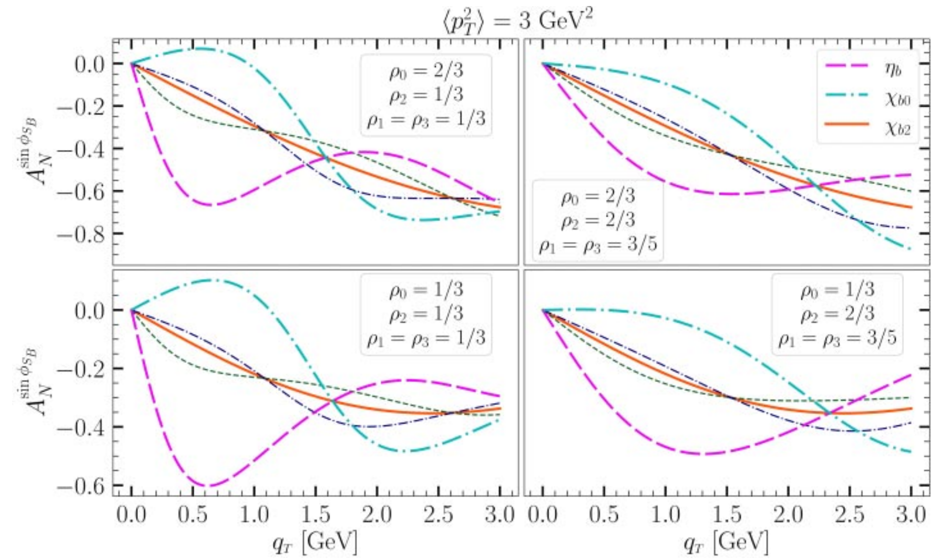

Fig. 4. – Same as in fig. 3, but for $\langle p_T^2 \rangle = 3$ GeV^2.

predictions of the maximal values of transverse single-spin asymmetries, by employing Gaussian parametrization for the gluon TMDs. The observables considered are sensitive to the gluon Sivers function $f_{1T}^{\perp g}$ in the case of $\chi_{\mathcal{Q}2}$ production. Once the Sivers function is determined, it can be used as input to isolate the effects of the linearly polarized gluon TMDs, h_1^g and $h_{1T}^{\perp g}$, which contribute to the SSAs for $\eta_{\mathcal{Q}}$ and $\chi_{\mathcal{Q}0}$ production. Such measurements should, in principle, be possible at LHCSpin, the fixed-target experiment planned at the LHC.

* * *

The author wishes to acknowledge L. Maxia and C. Pisano for their collaboration on this topic.

REFERENCES

[1] Bodwin G. T., Braaten E. and Lepage G. P., *Phys. Rev. D*, **51** (1995) 1125.
[2] Baier R. and Ruckl R., *Z. Phys. C*, **19** (1983) 251.
[3] LHCSpin Collaboration (Aidala A. C. *et al.*), *PoS*, **DIS2019** (2019) 233.
[4] LHCSpin Collaboration (Accardi A. *et al.*), arXiv:2504.16034.
[5] Boer D. and Pisano C., *Phys. Rev. D.*, **86** (2012) 094007.
[6] Kato N., Maxia L. and Pisano C., *Phys. Rev. D*, **110** (2024) 034038.
[7] Boer D., den Dunnen W. J., Pisano C., Schlegel M. and Vogelsang W., *Phys. Rev. Lett.*, **108** (2012) 032002.
[8] Lansberg J. P., Pisano C., Scarpa F. and Schlegel M., *Phys. Lett. B*, **784** (2018) 217.
[9] Mulders P. J. and Rodrigues J., *Phys. Rev. D*, **63** (2001) 094021.
[10] D'Alesio U., Murgia F., Pisano C. and Taels P., *Phys. Rev. D*, **96** (2017) 036011.
[11] D'Alesio U., Murgia F., Pisano C. and Rajesh S., *Eur. Phys. J. C*, **79** (2019) 1029.
[12] D'Alesio U., Maxia L., Murgia F., Pisano C. and Rajesh S., *Phys. Rev. D*, **102** (2020) 094011.