

Spin-spin entanglement in diffractive heavy-quark production

Michael Fucilla¹ and Yoshitaka Hatta^{2,3}

¹*National Centre for Nuclear Research, Pasteura 7, Warsaw 02-093, Poland*

²*Physics Department, Brookhaven National Laboratory, Upton, New York 11973, USA*

³*RIKEN BNL Research Center, Brookhaven National Laboratory, Upton, New York 11973, USA*



(Received 12 September 2025; accepted 3 February 2026; published 19 February 2026)

We calculate the spin density matrix of a heavy-quark-antiquark pair ($b\bar{b}$, $c\bar{c}$ or $s\bar{s}$) diffractively produced in deep inelastic scattering and ultraperipheral collisions. We show that the Pomeron exchange leaves characteristic imprints on the entanglement pattern between the quark and the antiquark. For the longitudinally polarized virtual photon, the pair always exhibits maximal entanglement and maximal violation of the Bell-Clauser-Horne-Shimony-Holt inequality. For the transversely polarized photon, the pair is always entangled and Bell violating, reaching maximal entanglement and maximal violation simultaneously when the transverse momentum approximately equals the quark mass.

DOI: [10.1103/gbk8-z3dd](https://doi.org/10.1103/gbk8-z3dd)

Introduction. The recent experimental measurements of the spin-spin entanglement in top-antitop ($t\bar{t}$) quark pairs by the ATLAS [1] and the CMS [2] Collaborations at the LHC have garnered significant attention amid a surge of interest in the intersection between collider physics and quantum information science [3–5]. While quantum entanglement among elementary particles such as electrons and photons has been studied since the early days of quantum mechanics [6–10], the direct observation of spin entanglement for quarks in QCD is highly nontrivial due to confinement. The top quarks are special because they decay by the electroweak interaction $t \rightarrow b + W^+$, $\bar{t} \rightarrow \bar{b} + W^-$ before the strong interaction kicks in. The information about the spin state of the $t\bar{t}$ pair can be recovered by measuring the angular distribution of the leptons $\ell^\pm$ produced in the subsequent decays $W^+ \rightarrow \ell^+ + \nu_\ell$, $W^- \rightarrow \ell^- + \bar{\nu}_\ell$ [11–15].

Another related paradigm is the test of Bell’s inequality [9] or, more specifically, the Clauser-Horne-Shimony-Holt (CHSH) inequality [10]. Bipartite states that violate the Bell-CHSH inequality are said to exhibit “Bell nonlocality,” a phenomenon that reflects more subtle and intricate quantum correlations than entanglement. In general, Bell nonlocality means entanglement, but the reverse is not true. As a result, establishing Bell nonlocality usually requires much stricter experimental tests [16,17] and has not yet been achieved in the top quark sector.

These quantum phenomena of course exist also in lighter quark systems such as bottom ($b\bar{b}$), charm ($c\bar{c}$), and strangeness ($s\bar{s}$) pairs. However, the corresponding measurements are more challenging because the pair first fragments hadronically. Fortunately, the spin state of the pair is partially retained when fragmenting into heavy baryons $q \rightarrow \Lambda_q$ and $\bar{q} \rightarrow \bar{\Lambda}_q$ [18,19] which then decay semileptonically. (See also an alternative approach [20].) Although such processes suffer from small branching ratios, measurements may still be possible [21,22] given the high luminosity of the LHC. The recent $\Lambda\bar{\Lambda}$ measurement at RHIC [23] is also encouraging.

It is natural to consider initiating similar experimental programs in deep inelastic scattering (DIS) at the Electron-Ion Collider (EIC) [24] that offers high luminosity and a cleaner experimental environment due to lepton scattering. While top quarks cannot be produced at the EIC, at least $c\bar{c}$ and $s\bar{s}$ pairs are copiously produced, and $b\bar{b}$ events can also be measured depending on kinematics. A related process is ultraperipheral collisions (UPCs) [25] where a heavy nucleus acts as a source of real photons. This mimics the photoproduction limit of DIS and can be studied already at the existing experimental facilities such as RHIC and the LHC. The spin density matrix of the $q\bar{q}$ pair from the lowest order process $\gamma^* + g \rightarrow q + \bar{q}$ has been recently calculated [26]. Remarkably, for the longitudinally polarized virtual photon γ_L^* , the produced $q\bar{q}$ pair is found to be always maximally entangled, and at the same time, the Bell-CHSH inequality is maximally violated.

The $2 \rightarrow 2$ process $\gamma^* + g \rightarrow q + \bar{q}$ corresponds to inclusive production where the target proton/nucleus breaks up. At the EIC, one can also study the exclusive diffractive production of $q\bar{q}$ pairs where the target stays intact.

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article’s title, journal citation, and DOI. Funded by SCOAP³.

This arises from the underlying process $\gamma^* + \mathbb{P} \rightarrow q + \bar{q}$ where $\mathbb{P}$ denotes “Pomeron,” the color-singlet gluonic exchange in the t -channel. Despite decades of intense scrutiny, the nature of the QCD Pomeron and the associated high density gluonic matter is not fully understood. In this Letter, we investigate the quantum informational aspect of the Pomeron by calculating the spin density matrix in $\gamma^* + \mathbb{P} \rightarrow q + \bar{q}$. We will be particularly interested in to what extent the Pomeron entangles the $q\bar{q}$ pair in spin space and whether the pair observes the Bell-CHSH inequality. Our finding provides theoretical motivation for a new interdisciplinary research direction that can be explored at the existing and future facilities. For previous works on different types of entanglement in the context of high energy QCD, see [27–39].

Heavy-quark pair production in diffractive DIS. Consider exclusive heavy-quark pair production in unpolarized electron-proton (or electron-nucleus) scattering $e + p \rightarrow e' + \gamma^* + p \rightarrow e' + q + \bar{q} + p'$.

We have in mind bottom, charm, and strange quarks $q = b, c, s$. First we work in a frame in which the virtual photon with momentum q^μ and virtuality $q^2 = -Q^2$ moves fast in the $+x^3$ direction and a $q\bar{q}$ pair is created in the photon fragmentation region with momenta

$$\tilde{k}^\mu = \left(zq^+, \frac{k_\perp^2 + m^2}{2zq^+}, \mathbf{k} \right), \quad \tilde{k}'^\mu = \left(\bar{z}q^+, \frac{k_\perp^2 + m^2}{2\bar{z}q^+}, -\mathbf{k} \right),$$

where m is the quark mass. z ($\bar{z} = 1 - z$) is the momentum fraction of the photon carried by the (anti)quark. It is related to the rapidities $y, \bar{y}$ as $z = e^y/(e^y + e^{\bar{y}})$. The light-cone coordinates $p^\mu = (p^+, p^-, \mathbf{p})$ are defined by $p^\pm = \frac{1}{\sqrt{2}}(p^0 \pm p^3)$ and two-dimensional transverse vectors are denoted by boldface letters $\mathbf{p} = (p^1, p^2)$ with $p_\perp \equiv |\mathbf{p}|$. The pair has invariant mass

$$M^2 = (\tilde{k} + \tilde{k}')^2 = \frac{k_\perp^2 + m^2}{z\bar{z}}. \quad (1)$$

For simplicity, the proton recoil momentum $\Delta = -\mathbf{k} - \mathbf{k}' \approx 0$ has been neglected.

We assume the “Regge” kinematics where the $\gamma^* + p$ center-of-mass (c.m.) energy $W^2 = (p + q)^2$ is large. In the eikonal approximation, for the longitudinally polarized virtual photon, the cross section is given by [40–44],

$$\left. \frac{d\sigma_{\alpha\alpha'\beta\beta'}^L}{dzd^2\mathbf{k}d^2\Delta} \right|_{\Delta \approx 0} = \frac{\alpha_{\text{em}} e_q^2 z\bar{z} Q^2}{N_c (q^+)^2} \left| \int \frac{d^2\mathbf{p} T(\mathbf{p})}{(\mathbf{k} - \mathbf{p})^2 + \mu^2} \right|^2 \times \bar{v}_{\beta'}(\tilde{k}') \gamma^+ u_\beta(\tilde{k}) \bar{u}_\alpha(\tilde{k}) \gamma^+ v_{\alpha'}(\tilde{k}'), \quad (2)$$

where $\mu^2 \equiv z\bar{z}Q^2 + m^2$ and e_q is the quark electromagnetic charge in units of $|e|$. α, α' and β, β' denote the spin states of the $q\bar{q}$ pair in the amplitude and the complex-conjugate

amplitude, respectively. Since our goal is to compute the spin density matrix, these indices have not been summed over [11,45,46]. The T -matrix of the $q\bar{q}$ pair (“color dipole”) T represents the Pomeron exchange in the present context. Our normalization is such that, in the Golec-Biernat–Wusthoff (GBW) model [47],

$$T(\mathbf{p}) = \frac{N_c \sigma_0}{(2\pi)^2} \left(\delta^{(2)}(\mathbf{p}) - \frac{R^2}{\pi} e^{-R^2 p_\perp^2} \right), \quad (3)$$

where σ_0 is the effective transverse area of the proton and R depends on W and Q [47]. In the following, we use the notations

$$T_1 \equiv \int \frac{d^2\mathbf{p} T(\mathbf{p})}{(\mathbf{k} - \mathbf{p})^2 + \mu^2}, \quad T_2 \equiv \frac{-1}{k_\perp^2} \int \frac{d^2\mathbf{p} \mathbf{k} \cdot \mathbf{p} T(\mathbf{p})}{(\mathbf{k} - \mathbf{p})^2 + \mu^2}. \quad (4)$$

To evaluate (2), we first go to the pair’s c.m. frame via a Lorentz boost $\tilde{k}^\pm = e^{\pm\eta} k^\pm$, $\tilde{k}'^\pm = e^{\pm\eta} k'^\pm$ with the boost factor $e^\eta = q^+ \sqrt{\frac{2z\bar{z}}{k_\perp^2 + m^2}}$. In this frame, the quark has velocity

$$|\vec{k}| = \frac{M}{2} \beta, \quad \beta = \sqrt{1 - \frac{4m^2}{M^2}} \quad (5)$$

and propagates in the direction

$$\cos \theta = \frac{k^3}{|\vec{k}|} = \frac{(z - \bar{z})M}{\sqrt{M^2 - 4m^2}}, \quad \phi = \arg(k^1 + ik^2), \quad (6)$$

in the polar coordinates measured from the $+x^3$ axis. The resulting spinor bilinear $\bar{u}(k)\gamma^+v(k')$ is most conveniently analyzed in the coordinate system spanned by the orthonormal basis $\{\hat{n}, \hat{r}, \hat{k}\}$ [11,12] where $\hat{k} = (\sin \theta \cos \phi, \sin \theta \sin \phi, \cos \theta)$ and

$$\hat{n} = \frac{\hat{x}^3 \times \hat{k}}{\sin \theta}, \quad \hat{r} = \frac{\hat{x}^3 - \hat{k} \cos \theta}{\sin \theta}. \quad (7)$$

We can then use the helicity basis along the $\hat{k}$ direction [48] and express (2) in the form

$$\frac{d\sigma_{\alpha\alpha'\beta\beta'}^L}{dzd^2\mathbf{k}d^2\Delta} = \frac{A^L}{4} (\delta_{\alpha\beta} \delta_{\alpha'\beta'} + C_{ab}^L \xi_a^\dagger \sigma^a \xi_\beta \eta_{\alpha'}^\dagger \sigma^b \eta_{\beta'}), \quad (8)$$

where $a, b = n, r, k$, and σ^a are the Pauli matrices. ξ, η are the two-component quark and antiquark spinors, respectively. We find

$$A^L = \frac{8\alpha_{\text{em}} e_q^2 z^2 \bar{z}^2 Q^2}{N_c} T_1^2, \quad (9)$$

and

$$\begin{aligned}
C_{nn}^L &= 1, \\
C_{rr}^L &= -C_{kk}^L = -\frac{1 - (2 - \beta^2)\cos^2\theta}{1 - \beta^2\cos^2\theta}, \\
C_{rk}^L &= C_{kr}^L = -\frac{\sqrt{1 - \beta^2}\sin 2\theta}{1 - \beta^2\cos^2\theta}.
\end{aligned} \quad (10)$$

Note that $(C_{rr}^L)^2 + (C_{rk}^L)^2 = 1$. The sign convention of C_{ab} is such that the matrix C_{ab} represents the correlation between the spin projections of the quark and the antiquark along the $+\hat{k}$ axis. For the antiquark moving in the $-\hat{k}$ direction, this is opposite to the helicity. It turns out that the matrix elements C_{ab}^L do not depend on Q and are exactly the same as for the “one-gluon exchange” process $\gamma_L^* + g \rightarrow q + \bar{q}$ [26]. This is somewhat surprising because the Pomeron exchanged in the t -channel consists of an arbitrary number of gluons $\mathbb{P} = gg, ggg, \dots$ in an overall color-singlet state. It also means that C_{ab}^L carries no information about the property of the target proton. As observed in [26], the states represented by the density matrix (10) are maximally entangled pure states. In the relativistic limit $\beta \rightarrow 1$ and/or for the symmetric configuration $z = \frac{1}{2}$ (meaning $\cos\theta = 0$ in the c.m. frame), it reduces to one of the Bell states.

Another measure of quantum correlation is the violation of the Bell-CHSH inequality [10]

$$\text{Max}_{\{\vec{n}_i\}} \left| n_1^a C_{ab} (n_2^b + n_4^b) + n_3^a C_{ab} (n_2^b - n_4^b) \right| \leq 2, \quad (11)$$

for unit vectors $|\vec{n}_i| = 1$ ($i = 1, 2, 3, 4$). This inequality is violated if the largest two of the three eigenvalues $\mu_3 \leq \mu_2 \leq \mu_1$ of the matrix $C^T C$ (the symbol T denotes “transpose”) satisfy [49]

$$1 < \mu_1 + \mu_2 \leq 2. \quad (12)$$

In the longitudinal case, we find $(C^L)^T C^L = \text{diag}(1, 1, 1)$ so that $\mu_1 + \mu_2 = 2$ and the inequality is maximally violated, with the left-hand side of (11) reaching $2\sqrt{2}$. Therefore, the $q\bar{q}$ pair always exhibits maximal entanglement and maximal Bell nonlocality both in single- and multiple-gluon exchanges.

Transverse photon. The calculation is significantly more complicated for the transversely polarized virtual photon. Again we write the $\gamma_T^* + p$ cross section in the form (8) with $L \rightarrow T$. The same cross section in the photoproduction limit $Q = 0$ is relevant to proton-nucleus UPCs [50–53]. The unpolarized cross section is well known

$$A^T = \frac{2\alpha_{em}e_q^2}{N_c} ((z^2 + \bar{z}^2)k_\perp^2 (T_1 + T_2)^2 + m^2 T_1^2). \quad (13)$$

The result for C_{ab}^T is new but lengthy, and we present it in Supplemental Material [54] together with a brief outline of

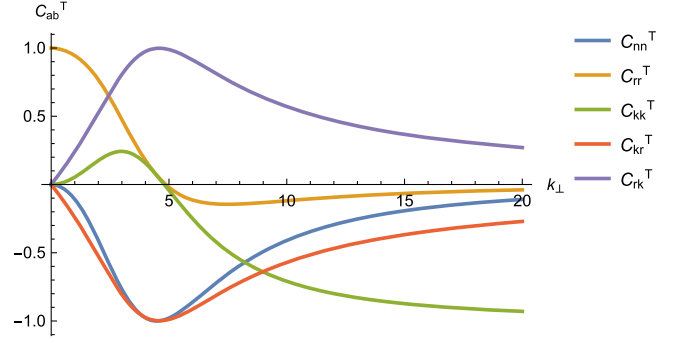


FIG. 1. C^T -matrix at $z = 1/2$ as a function of $k_\perp$ (in units of GeV) for the bottom quark $m = 4.18$ GeV. We used the GBW model [47] with $Q^2 = 9$ GeV² and $W = 100$ GeV.

the calculation. In contrast to the longitudinal case (10), C_{ab}^T is neither symmetric nor antisymmetric. (However, $C_{kr}^T = -C_{rk}^T$ for $z = \frac{1}{2}$.) Yet, we noticed the following nontrivial identities for generic values of z :

$$(C_{rr}^T)^2 + (C_{rk}^T)^2 + (C_{kr}^T)^2 + (C_{kk}^T)^2 - (C_{nn}^T)^2 = 1. \quad (14)$$

$$C_{nn}^T = -C_{rr}^T C_{kk}^T + C_{rk}^T C_{kr}^T. \quad (15)$$

In fact, the same identities hold in the longitudinal case (10) but in a more trivial manner. Figure 1 illustrates the behavior of the coefficients C_{ab}^T as a function of $k_\perp$ at $z = \frac{1}{2}$. Each curve shows a characteristic behavior around $k_\perp \approx \mu$.

In contrast to the one-gluon exchange studied in [26], C_{ab}^T now depends on the structure of the target proton through the dipole T-matrix $T_{1,2}$. Remarkably, however, the criterion for entanglement does not depend on $T_{1,2}$. Consider the following two quantities:

$$\begin{aligned}
\Delta_1 &= \sqrt{(C_{rr}^T - C_{kk}^T)^2 + (C_{rk}^T + C_{kr}^T)^2} - 1 + C_{nn}^T, \\
\Delta_2 &= \sqrt{(C_{rr}^T + C_{kk}^T)^2 + (C_{rk}^T - C_{kr}^T)^2} - 1 - C_{nn}^T.
\end{aligned} \quad (16)$$

According to the Peres-Horodecki criterion [56,57], if one of the Δ 's is non-negative, the $q\bar{q}$ pair is entangled [4,12]. It follows from (14) and (15) that

$$\Delta_2 = -\Delta_1 = -2C_{nn}^T \geq 0. \quad (17)$$

Therefore, the $q\bar{q}$ pair is always entangled except when $C_{nn}^T = 0$, i.e., on the kinematical end points $z = 0, 1$, or exactly at the threshold $k_\perp = 0$, or in the ultrarelativistic limit $\beta \rightarrow 1$ ($k_\perp \rightarrow \infty$). Again, this is in contrast to [26] (see also [22]) where there is a region of no entanglement (i.e., the system is separable), and entanglement grows stronger near the threshold or in the ultrarelativistic limit. Δ_2 is plotted as a function of z and $k_\perp$ in Fig. 2. For a fixed value of $k_\perp$, Δ_2 has a peak at $z = \frac{1}{2}$ where

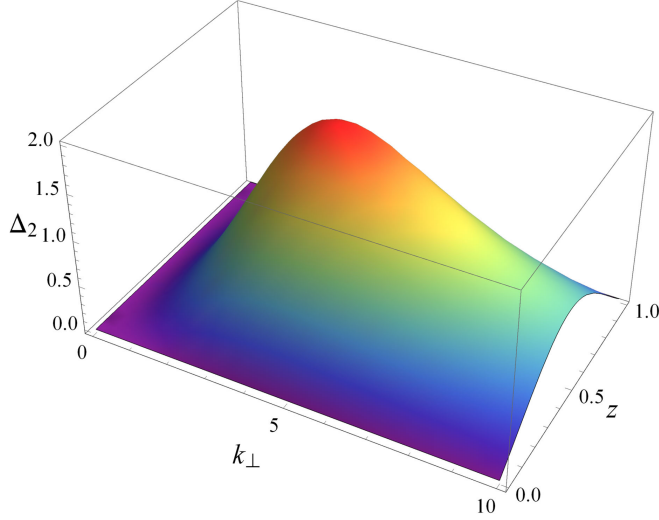


FIG. 2. Δ_2 in the $(z, k_\perp)$ plane at $Q^2 = 9 \text{ GeV}^2$, $m = 4.18 \text{ GeV}$, and $W = 100 \text{ GeV}$.

$$C_{nn}^T = -\frac{k_\perp^2 (T_1 + T_2)^2}{k_\perp^2 (T_1 + T_2)^2 + 2m^2 T_1^2}. \quad (18)$$

This takes the minimum value $C_{nn}^T = -1$ and $\Delta_2 = 2$ takes the maximal value if and only if $T_1 = 0$. [Naively, C_{nn}^T also approaches -1 in the ultrarelativistic limit $k_\perp \rightarrow \infty$. However, C_{nn}^T vanishes in this limit (cf. Fig. 1) due to a cancellation between T_1 and T_2 .] Interestingly, in heavy-quark production, T_1 does cross zero. This can be shown model independently by expanding the integrand of (4) in $|\mathbf{p}|$. Using the color transparency condition $\int d^2\mathbf{p} T(\mathbf{p}) = 0$, we find

$$T_1 \approx \frac{k_\perp^2 - \mu^2}{(k_\perp^2 + \mu^2)^3} \int d^2\mathbf{p} \mathbf{p}^2 T(\mathbf{p}). \quad (19)$$

Therefore, T_1 vanishes at $k_\perp \approx \mu = \sqrt{Q^2/4 + m^2}$, and at this point the C^T -matrix takes the form

$$C_{ab}^T \left(k_\perp \approx m, z = \frac{1}{2} \right) \approx \begin{pmatrix} -1 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & -1 & 0 \end{pmatrix}, \quad (20)$$

where we have taken the heavy-quark mass limit. The corresponding density matrix can be written as

$$\rho^T = \frac{1}{4} (1 \otimes 1 + C_{ab}^T \sigma^a \otimes \sigma^b) \approx |\Psi\rangle\langle\Psi|, \quad (21)$$

with

$$|\Psi\rangle = \frac{1}{\sqrt{2}} (|+\rangle_q^n |-\rangle_{\bar{q}}^n - i |-\rangle_q^n |+\rangle_{\bar{q}}^n). \quad (22)$$

This is a maximally entangled state, with the superscript n indicating that the spin quantization axis is taken along the $+\hat{n}$ direction (not the $+\hat{k}$ direction).

Next, we test the CHSH inequality. With the help of (14) and (15), it is easy to check that the three eigenvalues of the matrix $(C^T)^T C^T$ are

$$\mu_1 = 1, \quad \mu_2 = \mu_3 = (C_{nn}^T)^2. \quad (23)$$

Therefore, the CHSH inequality is violated as long as $C_{nn}^T < 0$, which is remarkably the same condition as for nonvanishing entanglement. This coincidence can be attributed to the identities (14) and (15). In particular, when $C_{nn}^T = -1$ where $\mu_1 + \mu_2 = 2$, the inequality is maximally violated. We therefore find that, even in the transversely polarized case, entanglement necessarily entails Bell nonlocality and vice versa. This is in sharp contrast to the situation in the one-gluon exchange [26] or, more generally, in typical two-qubit problems where Bell nonlocality is more nontrivial to realize than entanglement.

According to Gisin's theorem [58], for any two-qubit system that is a pure entangled state, the Bell-CHSH inequality is violated. Indeed, in the longitudinally polarized case, the density matrix satisfies the pure state condition $(\rho^L)^2 = \rho^L$. However, the transverse density matrix ρ^T does not in general represent a pure state since $(\rho^T)^2 \neq \rho^T$. The equality holds only if $C_{nn} = -1$, in which case the state becomes pure and maximally entangled (22). This is depicted in Fig. 3 (bottom), to be compared with the nesting structure for generic two-qubit systems (top).

Finally, we briefly discuss prospects for measurements using UPCs as an example, while leaving detailed feasibility studies for future work. In order to access spin information, the following channels have been identified as promising [21]:

$$\begin{aligned} b &\xrightarrow{7.0\%} \Lambda_b \xrightarrow{11\%} X_c^+ + \mu^- + \bar{\nu}_\mu, \\ c &\xrightarrow{6.4\%} \Lambda_c^+ \xrightarrow{3.5\%} \Lambda + \mu^+ + \nu_\mu, \\ s &\xrightarrow{2.8\%} \Lambda \xrightarrow{64\%} p + \pi^-, \end{aligned} \quad (24)$$

and their charge-conjugate counterparts for $\bar{b}, \bar{c}, \bar{s}$. The numbers above the first arrow are the fragmentation probabilities. The spin of the (anti)quark is largely inherited by the daughter (anti)baryon [18]. Similarly, especially for heavy quarks, roughly $z^{\text{quark}} \approx z^{\text{hadron}}$ and $k_\perp^{\text{quark}} \approx k_\perp^{\text{hadron}}$, to first approximation. The numbers above the second arrow are the branching ratios. Once these events have been identified and the baryon and antibaryon momenta have been reconstructed from their decay products [18,21,22], one goes to their respective rest frames by first moving to the pair's c.m. frame and then boosting along the $\pm\hat{k}$ directions. In each rest frame, the angular distribution of one of the decay particles can be measured. In order to

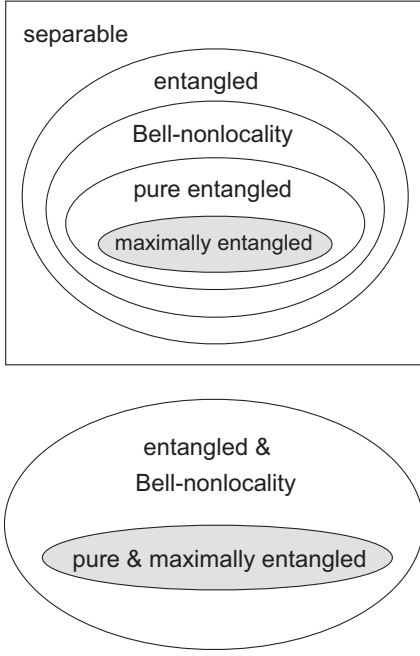


FIG. 3. Top: Venn diagram for generic two-qubit systems. Bottom: Venn diagram of a $q\bar{q}$ pair produced in $\gamma_T^* + \mathbb{P} \rightarrow q + \bar{q}$. The “separable” region is the boundary of the entangled region. If the photon is longitudinally polarized, only the shaded (“pure and maximal”) region is present.

access the component C_{nn}^T , one needs the polar angle distribution of each decay particle with respect to the $\hat{n}$ axis in the form

$$\frac{1}{\sigma} \frac{d\sigma^{\text{UPC}}}{d\cos\theta_+ d\cos\theta_-} = \frac{1 + \alpha_+ \alpha_- r_T^2 \cos\theta_+ \cos\theta_- C_{nn}^T}{4}, \quad (25)$$

where α_+ and $\alpha_- = -\alpha_+$ are the spin analyzing power of the measured particle and antiparticle, and r_T is the transverse polarization retention factor [18]. Equivalently,

$$C_{nn}^T = \frac{9}{\alpha_+ \alpha_- r_T^2} \langle \cos\theta_+ \cos\theta_- \rangle. \quad (26)$$

Since $C_{nn}^T \leq 0$, we predict that $\langle \cos\theta_+ \cos\theta_- \rangle \geq 0$, and this signifies both entanglement and Bell nonlocality, although the effect of hadronization should be carefully investigated. To give a rough idea about the expected yield, we consider UPCs with the lead nucleus ($A = 208$, $Z = 82$) at the LHC energy $\sqrt{s_{NN}} = 8.1$ TeV. According to [53], $d\sigma_{c\bar{c}}^{\text{UPC}}/dk_\perp \sim \mathcal{O}(100 \mu\text{b}/\text{GeV})$ at $k_\perp \sim m_c$. We have confirmed this result and moreover predict that $d\sigma_{b\bar{b}}^{\text{UPC}}/dk_\perp \sim \mathcal{O}(10 \text{ nb}/\text{GeV})$ at $k_\perp \sim m_b$ and $d\sigma_{s\bar{s}}^{\text{UPC}}/dk_\perp \sim \mathcal{O}(100 \mu\text{b}/\text{GeV})$. (Here we quote the value at $k_\perp \gtrsim 1$ GeV since, for the s -quark,

the region $k_\perp \sim m_s$ is nonperturbative.) The fragmentation and branching probabilities (24) bring in a suppression factor of order 10^{-4} for the $s\bar{s}$ pair and 10^{-6} for the $c\bar{c}$, $b\bar{b}$ pairs. Further considering $\alpha_\pm, r_T \sim \mathcal{O}(1)$ [18], we find the signal cross sections of order 10 nb, 100 pb, 0.01 fb for the $s\bar{s}$, $c\bar{c}$, $b\bar{b}$ channels, respectively. Assuming $1 \sim 10 \text{ pb}^{-1}$ of proton-lead collision data at the future high-luminosity LHC experiment [59], our rough estimates indicate that the $s\bar{s}$ -channel is feasible, while the $c\bar{c}$ -channel is marginal and the $b\bar{b}$ -channel is impractical at the LHC. In future, the discussion will be extended to light quarks [60] that have larger cross sections. Also, other methods to measure the spin correlation without relying on weak decays should be explored [20].

Conclusions. We have calculated the spin density matrix of a $q\bar{q}$ pair originating from the Pomeron, the color-singlet gluonic exchange in the t -channel that governs QCD amplitudes at high energy. In the longitudinally polarized case, we find maximal entanglement and maximal violation of the Bell-CHSH inequality. This agrees with the result from the one-gluon exchange in [26], although such an agreement is nontrivial. More strikingly, the transversely polarized photon in DIS or UPCs always leads to entanglement and Bell nonlocality in the produced $q\bar{q}$ pair. In other words, in the present system, entanglement is a necessary and sufficient condition for Bell nonlocality. This remarkable feature of the Pomeron, revealed in this Letter for the first time, to the best of our knowledge, is quite peculiar from the viewpoint of quantum information science (see Fig. 3) and certainly deserves further investigation.

Acknowledgments. We thank Bowen Xiao for stimulating discussions. M.F. thanks the EIC theory institute of Brookhaven National Laboratory, where this work was initiated, for hospitality. Y.H. is supported by the U.S. Department of Energy under Contract No. DE-SC0012704, by LDRD funds from Brookhaven Science Associates, and also by the framework of the Saturated Glue (SURGE) Topical Theory Collaboration. The work of M.F. is supported by the ULAM fellowship program of NAWA No. BNI/ULM/2024/1/00065 “Color glass condensate effective theory beyond the eikonal approximation.”

Data availability. The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

- [1] G. Aad *et al.* (ATLAS Collaboration), *Nature (London)* **633**, 542 (2024).
- [2] A. Hayrapetyan *et al.* (CMS Collaboration), *Rep. Prog. Phys.* **87**, 117801 (2024).
- [3] A. J. Barr, M. Fabbrihesi, R. Floreanini, E. Gabrielli, and L. Marzola, *Prog. Part. Nucl. Phys.* **139**, 104134 (2024).
- [4] Y. Afik and J. R. M. de Nova, *Quantum* **6**, 820 (2022).
- [5] Y. Afik *et al.*, *Eur. Phys. J. Plus* **140**, 855 (2025).
- [6] A. Einstein, B. Podolsky, and N. Rosen, *Phys. Rev.* **47**, 777 (1935).
- [7] C. S. Wu and I. Shaknov, *Phys. Rev.* **77**, 136 (1950).
- [8] D. Bohm, *Phys. Rev.* **85**, 166 (1952).
- [9] J. S. Bell, *Physics (Long Island City, N.Y.)* **1**, 195 (1964).
- [10] J. F. Clauser, M. A. Horne, A. Shimony, and R. A. Holt, *Phys. Rev. Lett.* **23**, 880 (1969).
- [11] M. Baumgart and B. Tweedie, *J. High Energy Phys.* **03** (2013) 117.
- [12] Y. Afik and J. R. M. de Nova, *Eur. Phys. J. Plus* **136**, 907 (2021).
- [13] J. A. Aguilar-Saavedra, *Phys. Rev. D* **108**, 076025 (2023).
- [14] T. Han, M. Low, and T. A. Wu, *J. High Energy Phys.* **07** (2024) 192.
- [15] Z. Dong, D. Gonçalves, K. Kong, and A. Navarro, *Phys. Rev. D* **109**, 115023 (2024).
- [16] M. Fabbrihesi, R. Floreanini, and G. Panizzo, *Phys. Rev. Lett.* **127**, 161801 (2021).
- [17] C. Severi, C. D. E. Boschi, F. Maltoni, and M. Sioli, *Eur. Phys. J. C* **82**, 285 (2022).
- [18] M. Galanti, A. Giammanco, Y. Grossman, Y. Kats, E. Stamou, and J. Zupan, *J. High Energy Phys.* **11** (2015) 067.
- [19] S.-J. Lin, M.-J. Liu, D. Y. Shao, and S.-Y. Wei, *J. High Energy Phys.* **11** (2025) 082.
- [20] K. Cheng and B. Yan, *Phys. Rev. Lett.* **135**, 011902 (2025).
- [21] Y. Kats and D. Uzan, *J. High Energy Phys.* **03** (2024) 063.
- [22] Y. Afik, Y. Kats, J. R. M. de Nova, A. Soffer, and D. Uzan, *Phys. Rev. D* **111**, L111902 (2025).
- [23] STAR Collaboration, *arXiv:2506.05499*.
- [24] R. Abdul Khalek *et al.*, *Nucl. Phys.* **A1026**, 122447 (2022).
- [25] A. J. Baltz *et al.*, *Phys. Rep.* **458**, 1 (2008).
- [26] W. Qi, Z. Guo, and B.-W. Xiao, *arXiv:2506.12889*.
- [27] A. Kovner and M. Lublinsky, *Phys. Rev. D* **92**, 034016 (2015).
- [28] R. Peschanski and S. Seki, *Phys. Lett. B* **758**, 89 (2016).
- [29] D. E. Kharzeev and E. M. Levin, *Phys. Rev. D* **95**, 114008 (2017).
- [30] Y. Liu and I. Zahed, *Phys. Rev. D* **100**, 046005 (2019).
- [31] G. S. Ramos and M. V. T. Machado, *Phys. Rev. D* **102**, 034019 (2020).
- [32] S. Bhattacharya, R. Boussarie, and Y. Hatta, *Phys. Lett. B* **859**, 139134 (2024).
- [33] Y. Guo, X. Liu, F. Yuan, and H. X. Zhu, *Research* **2025**, 0552 (2025).
- [34] J. D. Brandenburg, H. Duan, Z. Tu, R. Venugopalan, and Z. Xu, *Phys. Rev. Res.* **7**, 013131 (2025).
- [35] Y. Hatta and J. Montgomery, *Phys. Rev. D* **111**, 014024 (2025).
- [36] A. Dumitru and E. Kolbusz, *Phys. Rev. D* **111**, 114033 (2025).
- [37] S. Agrawal and R. Abir, *Phys. Lett. B* **868**, 139802 (2025).
- [38] M. Ouchen and A. Prygarin, *Phys. Rev. D* **112**, 094027 (2025).
- [39] M. Hentschinski, H. Jung, and K. Kutak, *arXiv:2509.03400*.
- [40] J. Bartels, H. Lotter, and M. Wüsthoff, *Phys. Lett. B* **379**, 239 (1996); **382**, 449(E) (1996).
- [41] N. N. Nikolaev, W. Schafer, B. G. Zakharov, and V. R. Zoller, *J. Exp. Theor. Phys.* **97**, 441 (2003).
- [42] T. Altinoluk, N. Armesto, G. Beuf, and A. H. Rezaeian, *Phys. Lett. B* **758**, 373 (2016).
- [43] Y. Hatta, B.-W. Xiao, and F. Yuan, *Phys. Rev. Lett.* **116**, 202301 (2016).
- [44] R. Boussarie, A. V. Grabovsky, L. Szymanowski, and S. Wallon, *Phys. Rev. D* **100**, 074020 (2019).
- [45] W. Bernreuther and A. Brandenburg, *Phys. Rev. D* **49**, 4481 (1994).
- [46] A. Brandenburg, M. Flesch, and P. Uwer, *Phys. Rev. D* **59**, 14001 (1999).
- [47] K. J. Golec-Biernat and M. Wüsthoff, *Phys. Rev. D* **59**, 014017 (1998).
- [48] M. Jacob and G. C. Wick, *Ann. Phys. (N.Y.)* **7**, 404 (1959).
- [49] R. Horodecki, P. Horodecki, and M. Horodecki, *Phys. Lett. A* **200**, 340 (1995).
- [50] Y. Hagiwara, Y. Hatta, R. Pasechnik, M. Tasevsky, and O. Teryaev, *Phys. Rev. D* **96**, 034009 (2017).
- [51] M. Reinke Pelicer, E. Gräve De Oliveira, and R. Pasechnik, *Phys. Rev. D* **99**, 034016 (2019).
- [52] V. P. Gonçalves, G. Sampaio dos Santos, and C. R. Sena, *Nucl. Phys.* **A1000**, 121862 (2020).
- [53] B. Linek, A. Łuszczak, M. Łuszczak, R. Pasechnik, W. Schäfer, and A. Szczurek, *J. High Energy Phys.* **10** (2023) 179.
- [54] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/gbk8-z3dd> for the derivation of the spin density matrix in the transverse case, which includes Ref. [55].
- [55] M. E. Peskin and D. V. Schroeder, *An Introduction to Quantum Field Theory* (Addison-Wesley, Reading, MA, 1995).
- [56] A. Peres, *Phys. Rev. Lett.* **77**, 1413 (1996).
- [57] P. Horodecki, *Phys. Lett. A* **232**, 333 (1997).
- [58] N. Gisin, *Phys. Lett. A* **154**, 201 (1991).
- [59] D. d'Enterria *et al.*, *J. Phys. G* **52**, 090501 (2025).
- [60] Y. Hatta and J. Schoenleber, *arXiv:2511.04537*.