

Quantum steering and discord in hyperon-antihyperon system in electron-positron annihilation

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Hyperon-antihyperon pairs produced from electron-positron annihilations provide a promising platform for exploring quantum-information properties in particle physics. In this work, we analyze quantum steering and quantum discord in hyperon-antihyperon systems produced from e^+e^- annihilations, based on the X-shaped spin density matrix reconstructed via quantum state tomography. We show that the behaviors of these correlations differ from those observed in elementary particle-antiparticle systems such as top quarks and τ leptons, owing to the polarization effects specific to hyperon decays. From the results, we examine the hierarchy of quantum correlations in hyperon-antihyperon systems, Bell nonlocality $\subset$ steering $\subset$ entanglement $\subset$ discord. We also discuss the role of experimental loopholes and the impact of quantum decoherence in collider-based studies of quantum correlations.

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I. INTRODUCTION

In recent years, testing quantum nonlocal correlations at colliders has attracted a lot of interest in high-energy physics (HEP). Historically, experimental investigations of fundamental properties in quantum mechanics were primarily conducted in photonic and atomic systems [1,2]. Nowadays, with the progress in particle detection technology and accumulation of huge amount of experimental data, there has been a growing interest in studying quantum information properties at smaller length and higher energy scales [3].

Generally, the spin states of unstable particles produced in high-energy experiments can be extracted based on the quantum state tomography [4]. For example, as the heaviest fundamental particle we have known, the top quark serve as an ideal testing ground for the study of quantum information in high-energy physics [5–7]. The experimental

evidence for the entanglement in $t\bar{t}$ has been identified at the LHC [8]. Apart from top quarks, investigation of other particle systems at colliders such as τ leptons [9–11], massive gauge bosons [12–14], hyperons [15–19], other quarks [20], and vector mesons [21], have been carried out as well. Particularly, the quantum correlations in the hyperon-antihyperon systems produced from charmonium decays were proposed and can be tested at the Beijing Spectrometer III (BESIII) [17,18,22,23].

As an asymmetric form of the quantum correlation, the quantum steering or Einstein-Podolsky-Rosen (EPR) steering was first introduced by Schrödinger [24]. It corresponds to the ability that the measurement of one subsystem can influence the state of another correlated subsystem, even when the subsystems are spatially separated. Steerability is regarded as a kind of quantum nonlocality that is less restrictive than the Bell nonlocality [25,26].

The quantum discord represents another type of the quantum correlation [27,28]. It captures quantum nature of the correlation that cannot be explained purely by classical physics, even when there is no entanglement between the subsystems. It reflects a more subtle aspect of quantum mechanics, especially in the context of quantum computation and quantum communication [29–31]. Nowadays, the quantum steering and discord have been introduced into particle physics such as the framework of the neutrino oscillation [32,33] and top quark systems [34,35].

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TABLE I. Some parameters in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$, where $Y\bar{Y}$ is a pair of ground-state octet hyperons.

	Branching ratio ($\times 10^{-4}$)	α_ψ	$\Delta\Phi/\text{rad}$	Ref.
$J/\psi \rightarrow \Lambda\bar{\Lambda}$	$19.43 \pm 0.03 \pm 0.33$	$0.4748 \pm 0.0022 \pm 0.0031$	$0.7521 \pm 0.0042 \pm 0.0066$	[40,41]
$J/\psi \rightarrow \Sigma^+\bar{\Sigma}^-$	$10.61 \pm 0.04 \pm 0.36$	$-0.508 \pm 0.006 \pm 0.004$	$-0.270 \pm 0.012 \pm 0.009$	[42,43]
$J/\psi \rightarrow \Sigma^0\bar{\Sigma}^0$	$11.64 \pm 0.04 \pm 0.23$	$-0.4133 \pm 0.0035 \pm 0.0077$	$-0.0828 \pm 0.00 \pm 0.0033$	[44]
$J/\psi \rightarrow \Xi^-\bar{\Xi}^+$	$10.40 \pm 0.06 \pm 0.74$	$0.586 \pm 0.012 \pm 0.010$	$1.213 \pm 0.046 \pm 0.016$	[45,46]
$J/\psi \rightarrow \Xi^0\bar{\Xi}^0$	$11.65 \pm 0.04 \pm 0.43$	$0.514 \pm 0.006 \pm 0.0015$	$1.168 \pm 0.019 \pm 0.018$	[47,48]

In this work, we focus on the hyperon-antihyperon systems produced in $e^+e^- \rightarrow \gamma^*/\psi \rightarrow Y\bar{Y}$ and study the quantum steering and discord in the $Y\bar{Y}$ system from a quantum information perspective. Our investigation is mainly based on the two-qubit spin density matrix of the hyperon-antihyperon pairs. We present the steerability using the three-setting measurement inequality [36,37], and derive an analytical expression for the quantum discord from its original definition [28]. Combining the present results with our earlier analysis of Bell nonlocality and entanglement [17], we demonstrate that the hierarchy of quantum correlations is reflected in the corresponding nonzero regions. Besides, we examine how the timelike electromagnetic form factors (EMFFs) influence the structure of hyperon's quantum correlations, leading to behaviors distinct from those observed in elementary particle systems. Finally, we further consider two experimental aspects relevant to collider-based analyses, namely the treatment of loopholes and the effects of detector-induced quantum decoherence.

The paper is organized as follows: in Sec. II, we introduce the two-qubit spin density matrix for hyperon-antihyperon in e^+e^- annihilation. We investigate the quantum steering and discord in Secs. III and IV, respectively. In Sec. V, a hierarchy of quantum correlations is presented. The role of hyperon's timelike electromagnetic form factors is discussed in Sec. VI. The nonlocality loophole and quantum decoherence are addressed in Secs. VIII and IX. In Sec. X, a summary of the main results and an outlook are presented.

II. PRELIMINARIES

We start this section with the spin density matrix for the hyperon and antihyperon [38,39], as a form of two-qubit system:

$$\rho_{Y\bar{Y}} = \frac{1}{4} \left(1 \otimes 1 + \mathbf{B}^+ \cdot \boldsymbol{\sigma} \otimes 1 + 1 \otimes \mathbf{B}^- \cdot \boldsymbol{\sigma} + \sum_{i,j} C_{ij} \sigma_i \otimes \sigma_j \right), \quad (1)$$

where $\boldsymbol{\sigma} = (\sigma_1, \sigma_2, \sigma_3)$ are Pauli matrices, $\mathbf{B}^+$ and $\mathbf{B}^-$ are the spin polarization vectors for the hyperon (Y) and antihyperon ($\bar{Y}$) respectively, and C_{ij} is the 3×3 spin

correlation matrix between them. In the analytical calculation with CP symmetry, the spin polarization vectors of the hyperon and antihyperon have the form

$$\mathbf{B}^+ = \mathbf{B}^- = \left(0, \frac{\beta_\psi \sin \vartheta \cos \vartheta}{1 + \alpha_\psi \cos^2 \vartheta}, 0 \right), \quad (2)$$

and the spin correlation matrix is

$$C_{ij} = \frac{1}{1 + \alpha_\psi \cos^2 \vartheta} \times \begin{bmatrix} \sin^2 \vartheta & 0 & \gamma_\psi \sin \vartheta \cos \vartheta \\ 0 & -\alpha_\psi \sin^2 \vartheta & 0 \\ \gamma_\psi \sin \vartheta \cos \vartheta & 0 & \alpha_\psi + \cos^2 \vartheta \end{bmatrix}, \quad (3)$$

where $\alpha_\psi \in [-1, 1]$ is the decay parameter, $\Delta\Phi \in (-\pi, \pi]$ is the relative phase, and ϑ is the scattering angle between the electron beam and the outgoing hyperon. β_ψ and γ_ψ are defined as

$$\begin{aligned} \beta_\psi &= \sqrt{1 - \alpha_\psi^2} \sin(\Delta\Phi), \\ \gamma_\psi &= \sqrt{1 - \alpha_\psi^2} \cos(\Delta\Phi). \end{aligned} \quad (4)$$

The two parameters are measured at BESIII, and we list some of the results in Table I. Here, $\mathbf{B}^\pm$ and C_{ij} are defined in both hyperon's and antihyperon's rest frames. After that, we choose three axes to be

$$\hat{\mathbf{y}} = \frac{\hat{\mathbf{p}}_Y \times \hat{\mathbf{p}}_e}{|\hat{\mathbf{p}}_Y \times \hat{\mathbf{p}}_e|}, \quad \hat{\mathbf{z}} = \hat{\mathbf{p}}_Y, \quad \hat{\mathbf{x}} = \hat{\mathbf{y}} \times \hat{\mathbf{z}}, \quad (5)$$

which applies to both the hyperon and antihyperon.

According to local unitary equivalence [17], the quantum correlation in a bipartite system keeps invariant under the local unitary transformation, so we can convert the two-qubit density operator to a standard X state, as follows:

$$\begin{aligned} \rho_{Y\bar{Y}}^X &= \frac{1}{4} \left(1 \otimes 1 + a\sigma_z \otimes 1 + 1 \otimes a\sigma_z \right. \\ &\quad \left. + \sum_i t_i \sigma_i \otimes \sigma_i \right), \end{aligned} \quad (6)$$

where

$$a = \frac{\beta_\psi \sin \vartheta \cos \vartheta}{1 + \alpha_\psi \cos^2 \vartheta},$$

$$t_{1,2} = \frac{1 + \alpha_\psi \pm \sqrt{(1 + \alpha_\psi \cos 2\vartheta)^2 - (\beta_\psi \sin 2\vartheta)^2}}{2(1 + \alpha_\psi \cos^2 \vartheta)},$$

$$t_3 = \frac{-\alpha_\psi \sin^2 \vartheta}{1 + \alpha_\psi \cos^2 \vartheta}. \quad (7)$$

Here, the X state can simplify the following calculations. Besides, we emphasize that the transformation from Eqs. (1) to (6) is a local unitary transformation performed in fixed spin quantization axes. Experimentally, the local unitary transformation to the X -state form is applied only to the globally reconstructed density matrix, rather than on an event-by-event basis. Therefore, no event-dependent rotations enter the experimental averaging procedure demonstrated in Ref. [17]. As a result, the reconstructed density matrix is obtained from an average over the full event sample, rather than from a subset of events that would lead to a fictitious state [49].

As is shown in Eq. (7), the X -form density matrix in Eq. (6) depends on two parameters $(\alpha_\psi, \Delta\Phi)$ as a function of ϑ . By explicitly expressing the Pauli matrices in Eq. (6) following Ref. [17], the X -state $\rho_{Y\bar{Y}}^X$ can be derived into the form

$$\rho_{Y\bar{Y}}^X = \frac{1}{4} \begin{bmatrix} 1 + 2a + t_3 & 0 & 0 & t_1 - t_2 \\ 0 & 1 - t_3 & t_1 + t_2 & 0 \\ 0 & t_1 + t_2 & 1 - t_3 & 0 \\ t_1 - t_2 & 0 & 0 & 1 - 2a + t_3 \end{bmatrix}, \quad (8)$$

where the name X state comes from its resemblance to the letter X .

III. QUANTUM STEERING

Quantum steering, also known as EPR steering, was first proposed in 1935 [24]. It was reformulated in the modern quantum information theory by the authors of Ref. [25] in 2007. Similar to the local-hidden-variable (LHV) hypothesis which was introduced in the Bell inequality, the local-hidden-state (LHS) model was put forward for the steerability of quantum systems.

In the context of a two-qubit state ρ_{AB} shared by Alice and Bob, the quantum steering means that Alice can “steer” Bob’s state if and only if (iff) the measurement outcomes of Alice and Bob exhibit correlation that violates the LHS model [25]. This violation suggests that the quantum system demonstrates a nonlocal property that cannot be

accounted for by any LHVs, which distinguishes the quantum steering from classical correlations.

A. Detecting steerability via steering inequality

Similar to the Bell nonlocality, the quantum steering involves the violation of certain inequalities (known as steering inequalities) that serve as its criterion. If these inequalities are violated, it indicates that the system exhibits a genuine quantum behavior. Cavalcanti, Jones, Wiseman, and Reid developed an inequality known as the CJWR steering inequality to diagnose the steerability of a two-qubit state with three-setting measurements on each side [36,37],

$$F_3^{\text{CJWR}} \equiv \frac{1}{\sqrt{3}} \left| \sum_{k=1}^3 \text{Tr}[\rho(A_k \otimes B_k)] \right| \leq 1, \quad (9)$$

where $A_k = \mathbf{s}_k \cdot \boldsymbol{\sigma}$, $B_k = \mathbf{r}_k \cdot \boldsymbol{\sigma}$, $\mathbf{s}_k$ and $\mathbf{r}_k$ ($k = 1, 2, 3$) are unit vectors, and $\{\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3\}$ are three basis vectors in a Cartesian coordinate system (any two vectors are orthogonal).

By setting the unit vectors $\mathbf{r}_k$ as $\|\mathbf{C}\mathbf{r}_1\| = \|\mathbf{C}\mathbf{r}_2\| = \|\mathbf{C}\mathbf{r}_3\| = \sqrt{\text{Tr}(\mathbf{C}^T \mathbf{C})/3}$ and $\mathbf{s}_k = \mathbf{C}\mathbf{r}_k / \sqrt{\text{Tr}(\mathbf{C}^T \mathbf{C})/3}$, F_3^{CJWR} in Eq. (9) reaches its maximum [50], where C_{ij} is the 3×3 correlation matrix in Eq. (1). Then the consequent maximal violation reads

$$\mathcal{F}_3[\rho] \equiv \max_{\mathbf{s}_k, \mathbf{r}_k} F_3^{\text{CJWR}} = \sqrt{\text{Tr}(\mathbf{C}^T \mathbf{C})}. \quad (10)$$

By applying the inequality to the X -shape density matrix for $Y\bar{Y}$, we obtain the maximal violation of the CJWR steering inequality

$$\begin{aligned} \mathcal{F}_3[\rho_{Y\bar{Y}}] &= \sqrt{t_1^2 + t_2^2 + t_3^2} \\ &= \sqrt{1 + 2 \left(\frac{\alpha_\psi \sin^2 \vartheta}{1 + \alpha_\psi \cos^2 \vartheta} \right)^2 - 2 \left(\frac{\beta_\psi \sin \vartheta \cos \vartheta}{1 + \alpha_\psi \cos^2 \vartheta} \right)^2}, \end{aligned} \quad (11)$$

which is shown in Fig. 1 for a variety of hyperons. We can observe that $\mathcal{F}_3[\rho_{Y\bar{Y}}]$ is a function of $\cos \vartheta$ (ϑ is the scattering angle) and is symmetric with respect to $\vartheta = 90^\circ$, with the black horizontal line being the steering bound $\mathcal{F}_3 = 1$.

There is a clear evidence that all hyperon-antihyperon systems in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$ violate the CJWR steering inequality. Furthermore, the violation has a peak in the transverse scattering at $\vartheta = 90^\circ$ with the maximal value $\mathcal{F}_3 = \sqrt{1 + 2\alpha_\psi^2}$ which depends only on the decay parameter α_ψ . Therefore, the steerability for hyperon-antihyperon systems can be quantified through the violation of the

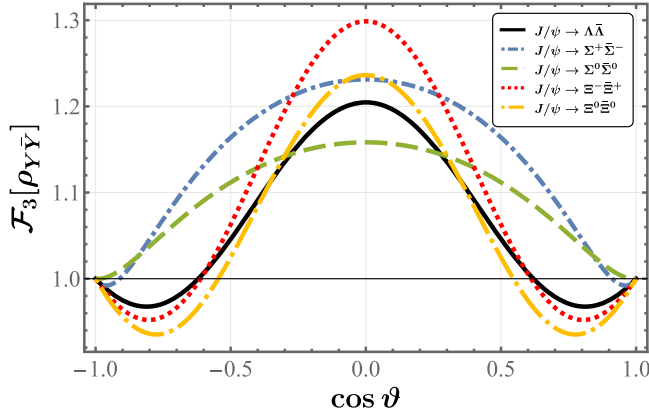


FIG. 1. Quantity $\mathcal{F}_3[\rho_{Y\bar{Y}}]$ for the three-setting measurement as functions of $\cos \vartheta$ (ϑ is the scattering angle) in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$ with the black solid, blue dash-dotted, green dashed, red dotted, and yellow long dash-dotted lines representing $Y = \Lambda$, Σ^+ , Σ^0 , Ξ^- , and Ξ^0 , respectively. The black horizontal line is the steering bound $\mathcal{F}_3 = 1$. The CJWR inequality is violated iff $\mathcal{F}_3 > 1$.

CJWR inequality. The ranges of steerability ($\vartheta^*, \pi - \vartheta^*$) for hyperon-antihyperon systems depend on both α_ψ and $\Delta\Phi$, where the critical angle ϑ^* is given by

$$\vartheta^* = \arctan \left| \sqrt{1 - \alpha_\psi^2} \frac{\sin \Delta\Phi}{\alpha_\psi} \right|. \quad (12)$$

The maximal violation for Σ is weaker than that for Ξ due to the relative small α_ψ , while the range of steerability for Σ is broader due to the very small value of $\Delta\Phi$.

It should be noted that the CJWR inequality serves only as a sufficient condition for the quantum steering. Consequently, states that do not violate the inequality are not necessarily unsteerable. Apart from the CJWR inequality, there are other criteria for the quantum steering as well [26]. For instance, an alternative approach to detect the quantum steering is presented in Appendix A through the entanglement [51,52]. However, most available criteria are sufficient but not necessary, limiting their ability to detect steerable states.

So far, the only known necessary and sufficient criterion for the quantum steering was introduced in Ref. [53] via the definition of the *critical radius*. A rigorous proof is given that a two-qubit state is steerable iff the critical radius is less than 1. This criterion has been applied to the top quark system [34]. However, a closed formula for the critical radius exists only for the Bell diagonal state, limiting its applicability to hyperon-antihyperon systems.

IV. QUANTUM DISCORD

The quantum discord is another type of the quantum property that quantifies how much quantum correlation shared between a bipartite system. Before discussing the

concept of quantum discord, we need to first introduce the *quantum mutual information* of a bipartite system ρ_{AB} , which has the form

$$\begin{aligned} \mathcal{I}(A:B) &\equiv S(\rho_A) + S(\rho_B) - S(\rho_{AB}) \\ &= S(\rho_A) - S(\rho_{A|B}), \end{aligned} \quad (13)$$

where $S(\rho) \equiv -\text{Tr}(\rho \log_2 \rho)$ denotes the von Neumann entropy, and $S(\rho_{A|B}) \equiv S(\rho_B) - S(\rho_{AB})$ is the conditional entropy.

To quantify the quantum discord, one can use a set of one-dimensional projectors proposed by Ollivier and Zurek [28]. Here, $\{\Pi_k, \sum_k \Pi_k = 1\}$ denotes a set of projectors that perform the projective measurement on subsystem B . The postmeasurement state of the subsystem A with the condition of the measurement outcome k of B reads

$$\rho_{A|\Pi_k} = \frac{1}{p_k} \text{Tr}_B[(1 \otimes \Pi_k) \rho_{AB} (1 \otimes \Pi_k)], \quad (14)$$

where the probability p_k is given by $\text{Tr}[(1 \otimes \Pi_k) \rho_{AB} (1 \otimes \Pi_k)]$. The postmeasurement states of A form an ensemble $\{p_k, \rho_{A|\Pi_k}\}$. Specifically, in a bipartite qubit system, $k = 0, 1$. Afterwards, the mutual information conditional on the projective measurement performed on the subsystem B is defined as

$$\mathcal{J}(A:B) \equiv S(\rho_A) - \sum_k p_k S(\rho_{A|\Pi_k}), \quad (15)$$

which is similar to Eq. (13). In order to avoid the measurement dependence, the classical mutual information of the state ρ_{AB} is given by maximizing the value of $\mathcal{J}(A:B)$ over all possible projectors.

The quantum discord is defined as the difference between the quantum and classical mutual information (13) and (15), as follows:

$$\begin{aligned} \mathcal{D}[\rho_{AB}] &\equiv \mathcal{I}(A:B) - \max_{\{\Pi_k\}} \mathcal{J}(A:B) \\ &= S(\rho_B) - S(\rho_{AB}) + \min_{\{\Pi_k\}} \sum_k p_k S(\rho_{A|\Pi_k}). \end{aligned} \quad (16)$$

This definition indicates that the quantum discord is an entropylike quantity, which is non-negative and can never exceed 1; i.e., $0 \leq \mathcal{D}[\rho_{AB}] \leq 1$. Another feature of the quantum discord is that it is an asymmetric quantity—performing the measurement on A or on B can return a different value. However, in our work, all hyperon-antihyperon states from $e^+e^- \rightarrow Y\bar{Y}$ are symmetric due to CP symmetry (6), ensuring identical values of $\mathcal{D}[\rho_{Y\bar{Y}}]$ regardless of whether the measurement is performed on Y or $\bar{Y}$.

A. Discord for rank-2X states

The obstacle to obtain the quantum discord lies in the sophisticated maximization (or minimization) procedure outlined in Eq. (16), which must be done over all possible projective measurements on B . The difficulty makes the exact expression of the quantum discord for general two-qubit states still a blank.

We know that the spin states of hyperon-antihyperon systems in e^+e^- scattering are *symmetric rank-2X states* with two nonzero eigenvalues [17], as follows:

$$\lambda_{1,2} = \frac{1}{2} \left(1 \mp \frac{\alpha_\psi \sin^2 \vartheta}{1 + \alpha_\psi \cos^2 \vartheta} \right). \quad (17)$$

Fortunately, there is an exact and analytical formula for the quantum discord of X states [54–57]:

$$\begin{aligned} \mathcal{D}[\rho_X] = & 1 - \frac{1+a}{2} \log_2 \frac{1+a}{2} - \frac{1-a}{2} \log_2 \frac{1-a}{2} \\ & + \sum_{i=1}^4 \lambda_i \log_2 \lambda_i - \max_{\epsilon \in [0,1]} F(\epsilon), \end{aligned} \quad (18)$$

where $F(\epsilon)$ is defined as

$$\begin{aligned} F(\epsilon) = & \frac{1+a\epsilon+H_+}{4} \log_2 \frac{1+a\epsilon+H_+}{1+a\epsilon} \\ & + \frac{1+a\epsilon-H_+}{4} \log_2 \frac{1+a\epsilon-H_+}{1+a\epsilon} \\ & + \frac{1-a\epsilon+H_-}{4} \log_2 \frac{1-a\epsilon+H_-}{1-a\epsilon} \\ & + \frac{1-a\epsilon-H_-}{4} \log_2 \frac{1-a\epsilon-H_-}{1-a\epsilon}, \end{aligned} \quad (19)$$

with $H_\pm = \sqrt{t^2(1-\epsilon^2) + (a \pm t_3\epsilon)^2}$, $t = \max\{|t_1|, |t_2|\}$. Equation (18) accounts for all two-qubit X states, and according to Ref. [56], the quantum discord of any rank-2 mixed states of X -type are always given by either $F(1)$ or $F(0)$. Simplifying Eq. (18) and using the numerical method, we find the discord for $Y\bar{Y}$ in our study can always be obtained at $\epsilon = 0$, so we have

$$\begin{aligned} \mathcal{D}[\rho_{Y\bar{Y}}] = & h\left(\frac{1+a}{2}\right) - h\left(\frac{1+t_3}{2}\right) \\ & + h\left(\frac{1+\sqrt{t_1+t_3-t_1t_3}}{2}\right), \end{aligned} \quad (20)$$

where

$$h(x) \equiv -x \log_2 x - (1-x) \log_2 (1-x), \quad (21)$$

is the Shannon binary entropy function, and a and $t_{1,2,3}$ are functions of ϑ which depend on parameters α_ψ and $\Delta\Phi$ as shown in Eq. (7).

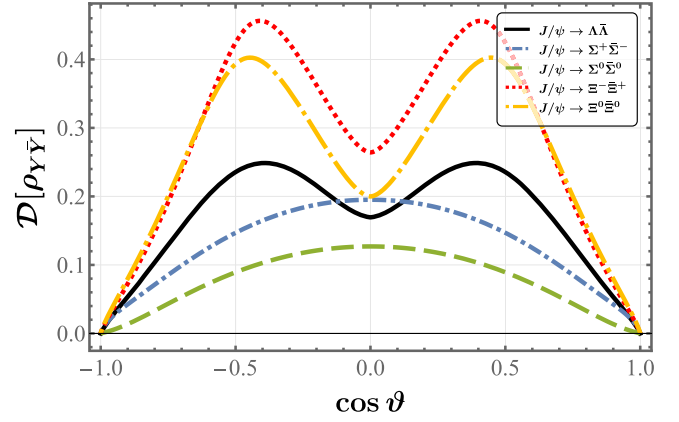


FIG. 2. Quantum discord $\mathcal{D}[\rho_{Y\bar{Y}}]$ as functions of $\cos \vartheta$ (ϑ is the scattering angle) in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$ with $Y = \Lambda, \Sigma^{+,0}$ and $\Xi^{-,0}$ corresponding to curves in black solid, blue dash-dotted, green dashed, red dotted, and yellow long dash-dotted lines, respectively. The black horizontal line is $\mathcal{D} = 0$.

The quantum discord $\mathcal{D}[\rho_{Y\bar{Y}}]$ is plotted as functions of $\cos \vartheta$ for $\Lambda, \Sigma^{+,0}$ and $\Xi^{0,-}$ hyperons in Fig. 2. We can observe that the discord of the hyperon-antihyperon system is symmetric with respect to $\vartheta = \pi/2$ (the transverse scattering angle). In the whole range of $\vartheta \in [0, \pi]$, the quantum discord of spin states of $Y\bar{Y}$ is nonzero. Additionally, unlike the steering, the maximum quantum discord is not necessarily located at $\vartheta = \pi/2$; for Λ and $\Xi^{-,0}$, the discord has two peaks away from $\vartheta = \pi/2$.

B. Geometric quantum discord

Since the quantum discord given in Eq. (16) is generally hard to evaluate, an alternative geometric measure for the discord is proposed from Ref. [58]:

$$\mathcal{D}_G \equiv \min_{\chi \in \Omega_0} \|\rho - \chi\|^2, \quad (22)$$

where Ω_0 denotes the set of zero-discord states, $\|\rho - \chi\|^2 = \text{Tr}[(\rho - \chi)^2]$ is the Hilbert-Schmidt distance between two states, and $\mathcal{D}_G$ is known as the *geometric quantum discord*.

The exact and analytical expression for two-qubit states has also been given in Ref. [58]. For the X state given in Eq. (6), the geometric discord can be evaluated as

$$\mathcal{D}_G[\rho_{Y\bar{Y}}] = \frac{1}{4} \min \{a^2 + t_2^2 + t_3^2, t_1^2 + t_2^2\}. \quad (23)$$

It is easy to see that $\mathcal{D}_G$ is not normalized to 1: its maximum value is 1/2 for two-qubit states, so it is natural to consider $2\mathcal{D}_G$ as a proper measure in comparison with the quantum discord $\mathcal{D}$ [59]. The results for $2\mathcal{D}_G$ are plotted in Fig. 3, which are nonzero in the full range of the scattering angle and agree with the original discord in Fig. 2. Since $\mathcal{D}_G$ is a geometric quantity while $\mathcal{D}$ is an entropylike quantity, the relationship between them needs further investigation.

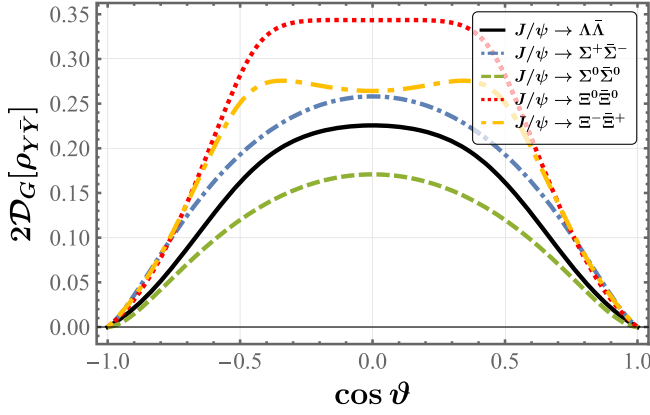


FIG. 3. Results for the geometric quantum discord $2\mathcal{D}_G[\rho_{Y\bar{Y}}]$ as functions of $\cos \vartheta$ (ϑ is the scattering angle) in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$ with $Y = \Lambda, \Sigma^+, \Sigma^0, \Xi^0$ and Ξ^- corresponding to black solid, blue dash-dotted, green dashed, red dotted, and yellow long dash-dotted lines, respectively. The black horizontal line is the zero geometric discord $2\mathcal{D}_G = 0$.

V. DIFFERENT TYPES OF QUANTUM CORRELATIONS

In this section, we present four types of quantum correlations within the hierarchy introduced in the framework of quantum information [25,34,60]:

$$\text{Bell nonlocality} \subset \text{steering} \subset \text{entanglement} \subset \text{discord}, \quad (24)$$

where the hierarchy refers to the inclusion relations among sets of quantum states. For hyperon-antihyperon systems, the Bell nonlocality and entanglement properties have been investigated in earlier works [17,18]. Here we bring together all four types of quantum correlations in a unified analysis. We show that, for the systems considered, the hierarchy in Eq. (24) is manifested through the nonzero regions of different correlations as functions of the scattering angle.

To be simplified, we scale measures of all the correlations to a range of $[0, 1]$. Since quantum discord is an entropylike quantity, we can use $\mathcal{D} \equiv \mathcal{D}[\rho_{Y\bar{Y}}]$ in Eq. (16). Although the concurrence is often employed as an entanglement measure, in this work we adopt *entanglement of formation* [61] with the definition of

$$\begin{aligned} \mathcal{E} &\equiv h\left(\frac{1 + \sqrt{1 - \mathcal{C}^2[\rho_{Y\bar{Y}}]}}{2}\right) \\ &= h\left(\frac{1 + \sqrt{1 - t_2^2}}{2}\right), \end{aligned} \quad (25)$$

where $h(x)$ is the Shannon binary entropy function, and $\mathcal{C}[\rho_{Y\bar{Y}}] = |t_2|$ in the second line [17]. Thus, the measure of entanglement can be taken as an entropylike quantity with $\mathcal{E} \in [0, 1]$ as well.

Bell nonlocality and quantum steering are quantified by the violation of the Bell and CJWR inequalities, respectively. The bound for the quantum steering is $\mathcal{F}_3 > 1$ and the maximal violation is $\mathcal{F}_3^{\max} = \sqrt{3}$. Therefore, we introduce a rescaled measure of the quantum steering, namely

$$\mathcal{S} \equiv \max\left\{0, \frac{\mathcal{F}_3 - 1}{\sqrt{3} - 1}\right\} \in [0, 1]. \quad (26)$$

Similarly, the bound for the Bell nonlocality is $\mathcal{B} > 2$ with the maximum violation to be $\mathcal{B}^{\max} = 2\sqrt{2}$. We can correspondingly rescale the measure to be

$$\mathcal{B} \equiv \max\left\{0, \frac{\mathcal{B} - 2}{2\sqrt{2} - 2}\right\} \in [0, 1]. \quad (27)$$

In Fig. 4, we compare four quantum correlations—Bell nonlocality (BN) $\mathcal{B}$, steering $\mathcal{S}$, entanglement $\mathcal{E}$ and discord $\mathcal{D}$ —as functions of $\cos \vartheta$ for various hyperon-antihyperon systems in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$, with panels (a)–(e) for $\Lambda, \Sigma^+, \Sigma^0, \Xi^0$, and Ξ^- , respectively. In each panel, $\mathcal{B}, \mathcal{S}, \mathcal{E}$, and $\mathcal{D}$ are plotted in the blue dash-dotted, yellow dashed, black solid, and red dotted curves, respectively.

As is presented in Fig. 4, the nonzero regions with respect to the scattering angle ϑ (or $\cos \vartheta$) provide a clear manifestation of the hierarchical relation in Eq. (24). In particular, both the discord and the entanglement are nonzero over the entire angular range, except at the two endpoints, $\vartheta = 0$ and 180° . In contrast, the steerability is restricted to a finite interval centered around the transverse scattering angle, while the Bell nonlocality is even more constrained and lies entirely within the steerable region.

Moreover, representative parameters characterizing the four quantum correlations are summarized in Table II. These parameters provide a more intuitive characterization of the hierarchy in Eq. (24) by quantifying the nonzero regions of different correlations. In particular, the table shows that steering is less restrictive than Bell nonlocality but more restrictive than entanglement, as reflected by their respective angular supports. By contrast, the quantum discord $\mathcal{D}$ remains nonzero over the entire scattering-angle range, similar to the entanglement.

In addition, we also calculate the quantum correlations in $Y\bar{Y}$ systems through $\psi(3686)$ in e^+e^- annihilation using the available data for the parameters α_ψ and $\Delta\Phi$; see Appendix B.

VI. EMFFS IN QUANTUM CORRELATIONS

In this section, we discuss the role of electromagnetic form factors in quantum correlations in hyperon-antihyperon systems. In $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$, the J/ψ -hyperon vertex can be written as [38]

$$\Gamma_\mu(p_1, p_2) = -ie \left[G_M \gamma_\mu - \frac{2M}{Q^2} (G_M - G_E) Q_\mu \right], \quad (28)$$

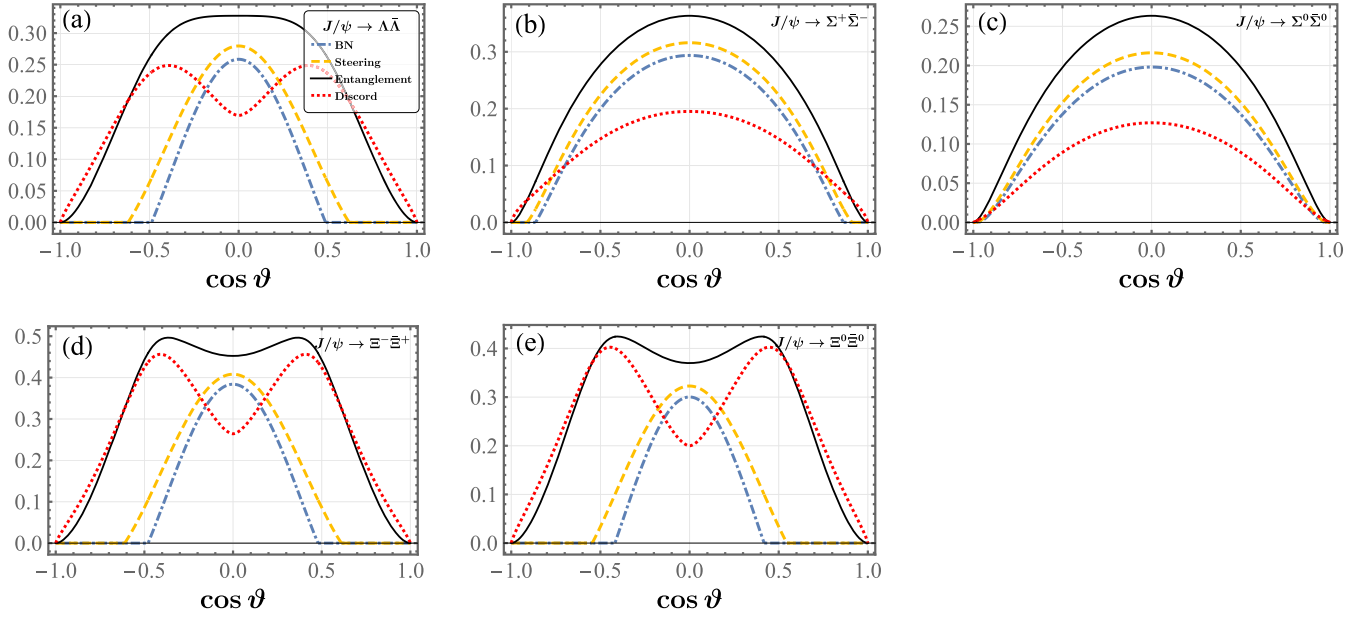


FIG. 4. Four different types of quantum correlations as functions $\cos \vartheta$ (ϑ is the scattering angle) in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$. The five panels from (a)–(e) correspond to Λ , Σ^+ , Σ^0 , Ξ^- , and Ξ^0 , respectively. The Bell nonlocality $\mathcal{B}$ is shown in blue dot-dashed lines, the steering $\mathcal{S}$ is shown in yellow dashed lines, the entanglement $\mathcal{E}$ is shown in black solid lines, and the discord $\mathcal{D}$ is shown in red dotted lines.

where $P = p_1 + p_2$, $Q = p_1 - p_2$, M is the hyperon mass, and G_E and G_M are electric and magnetic form factors of hyperons as functions $s = P^2$. The decay parameters α_ψ and $\Delta\Phi$ are defined through G_E and G_M by

$$\alpha_\psi = \frac{s|G_M|^2 - 4M^2|G_E|^2}{s|G_M|^2 + 4M^2|G_E|^2}, \quad G_E/G_M = e^{i\Delta\Phi}|G_E/G_M|, \quad (29)$$

where $\Delta\Phi$ is the relative phase between two complex-valued form factors.

Equation (7) shows that the quantum correlations at $\vartheta = 90^\circ$ only depend on the decay parameter α_ψ and are irrelevant of the phase $\Delta\Phi$, as follows:

$$\begin{aligned} \mathcal{D}_{\vartheta=90^\circ} &= 1 - h\left(\frac{1 - \alpha_\psi}{2}\right), \\ \mathcal{E}_{\vartheta=90^\circ} &= h\left(\frac{1 + \sqrt{1 - \alpha_\psi^2}}{2}\right), \\ \mathcal{S}_{\vartheta=90^\circ} &= \frac{\sqrt{1 + 2\alpha_\psi^2} - 1}{\sqrt{3} - 1}, \\ \mathcal{B}_{\vartheta=90^\circ} &= \frac{\sqrt{1 + \alpha_\psi^2} - 1}{\sqrt{2} - 1}. \end{aligned} \quad (30)$$

In order to study the effect of $\Delta\Phi$ in quantum correlations, we deliberately set $\Delta\Phi$ to zero. In this sense, the analytical expressions of these four types of correlations are reduced to simpler forms:

TABLE II. Maximal values and ranges for four types of quantum correlations. The angles ϑ_{BN}^* and $\vartheta_{\text{steering}}^*$ are listed for the nonzero Bell nonlocality and steering in the range $[\vartheta^*, 180^\circ - \vartheta^*]$. The angle ϑ^{max} shown in parentheses corresponds to the one at which the hyperon-antihyperon system reaches its maximum entanglement or discord.

	$\Lambda\bar{\Lambda}$	$\Sigma^+\bar{\Sigma}^-$	$\Sigma^0\bar{\Sigma}^0$	$\Xi^-\bar{\Xi}^+$	$\Xi^0\bar{\Xi}^0$
$\mathcal{B}_{\text{max}}$	0.259	0.294	0.198	0.384	0.300
ϑ_{BN}^*	27.93°	49.45°	14.30°	61.38°	65.29°
$\mathcal{S}_{\text{max}}$	0.280	0.316	0.216	0.408	0.323
$\vartheta_{\text{steering}}^*$	35.70°	52.08°	10.58°	52.41°	56.98°
$\mathcal{E}_{\text{max}}(\vartheta^{\text{max}})$	0.327(90°)	0.363(90°)	0.263(90°)	0.496(68.60°)	0.425(66.26°)
$\mathcal{D}_{\text{max}}(\vartheta^{\text{max}})$	0.170(90°)	0.195(90°)	0.127(90°)	0.456(65.86°)	0.403(63.49°)

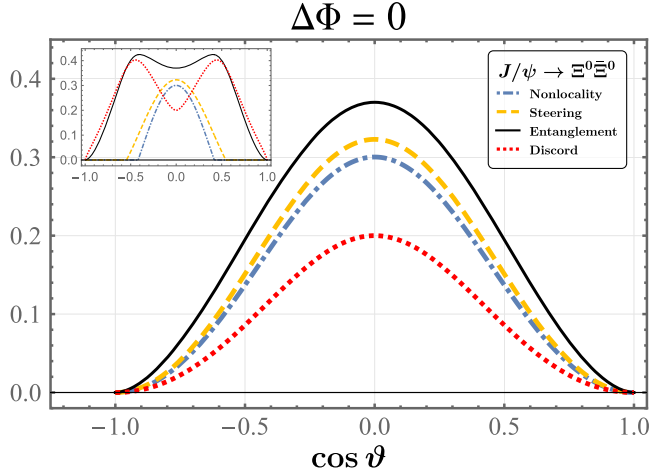


FIG. 5. Quantum correlations for $\Xi^0\bar{\Xi}^0$ with $\Delta\Phi = 0$. The subfigure shows the result in Fig. 4(e) for comparison. The Bell nonlocality $\mathcal{B}$, the steering $\mathcal{S}$, the entanglement $\mathcal{E}$, and the discord $\mathcal{D}$ are shown in blue dot-dashed, yellow dashed, black solid, and red dotted lines, respectively.

$$\begin{aligned}
 \mathcal{D}_{\Delta\Phi=0} &= 1 - h\left(\frac{1 + \alpha_\psi \cos 2\vartheta}{2(1 + \alpha_\psi \cos^2 \vartheta)}\right), \\
 \mathcal{E}_{\Delta\Phi=0} &= h\left(\frac{1}{2} + \frac{\sqrt{(1 + \alpha_\psi)(1 + \alpha_\psi \cos 2\vartheta)}}{2(1 + \alpha_\psi \cos^2 \vartheta)}\right), \\
 \mathcal{S}_{\Delta\Phi=0} &= \frac{\sqrt{1 + 2[\alpha_\psi \sin^2 \vartheta / (1 + \alpha_\psi \cos^2 \vartheta)]^2} - 1}{\sqrt{3} - 1}, \\
 \mathcal{B}_{\Delta\Phi=0} &= \frac{\sqrt{1 + [\alpha_\psi \sin^2 \vartheta / (1 + \alpha_\psi \cos^2 \vartheta)]^2} - 1}{\sqrt{2} - 1}. \quad (31)
 \end{aligned}$$

Here, we choose $J/\psi \rightarrow \Xi^0\bar{\Xi}^0$ as a representative to illustrate the effect of $\Delta\Phi$. We plot the correlations in Fig. 5 with and without $\Delta\Phi$. It shows that the profiles of four types of quantum correlations have been affected.

VII. QUANTUM CORRELATION AS PROBE TO PARTICLE COMPOSITENESS

As composite particles, hyperons exhibit spin-correlation patterns that are qualitatively different from those of elementary fermions. In particular, the nonzero relative phase between the EMFFs generates a polarization along the y direction, proportional to $\sin(\Delta\Phi)$ and thus to $\text{Im}(G_M G_E^*)$. Such correlations have no analog in elementary systems such as $t\bar{t}$ pairs [5–7,34] or $\tau^+\tau^-$ leptons [9,10]. The underlying difference arises from the interference between G_E and G_M in the J/ψ -hyperon vertex [Eq. (28)].

In contrast, $t\bar{t}$ production in $p\bar{p}$ collisions is through two partonic channels, $q\bar{q} \rightarrow t\bar{t}$ and $g\bar{g} \rightarrow t\bar{t}$. These channels yield a simple relation between Bell nonlocality and

concurrence, $\mathcal{B} = 2\sqrt{1 + \mathcal{C}^2}$, which also appears in $\tau^+\tau^-$ production [17].

This phenomenological distinction suggests a conjecture: the quantum correlation of particle-antiparticle may serve as a probe to the compositeness of the particle. The complex-valued structure of the hyperon's timelike EMFFs is therefore essential to understand the observed quantum correlation patterns. This motivates a deeper exploration of hyperon-antihyperon systems within the framework of quantum information at colliders such as Beijing Electron-Positron Collider II (BEPCII) and the proposed Super Tau-Charm Facility (STCF) [62].

VIII. THE LOCALITY LOOPHOLE IN HEP EXPERIMENTS

In the history of experimental tests of quantum nonlocality, several well-known loopholes have been identified, including the detection-efficiency loophole, the freedom-of-choice loophole, and the locality loophole [63,64]. If not properly taken into account, these loopholes may allow local realistic models to reproduce quantum predictions, thereby complicating the interpretation of experimental violations of Bell inequalities. Considerable experimental efforts have therefore been devoted to closing these loopholes in low-energy platforms, leading to the realization of loophole-free Bell tests in recent years [65,66].

In high-energy collider experiments, however, the situation is qualitatively different. Due to the nature of particle production, triggering, and decay, none of the standard loopholes can be closed in a device-independent manner. As a consequence, the interpretation and relevance of loopholes in collider-based studies of quantum correlations remain open and actively discussed topics in the literature [6,10,18,67–71].

In the following, we focus on the locality aspect, not because it is the only relevant loophole in collider experiments, but because it admits a clear and quantitative kinematic characterization within the framework of hyperon decays. The discussion below is therefore not intended to establish a loophole-free Bell test, but rather to illustrate how locality constraints enter the analysis of quantum correlations in collider environments.

As a concrete example, we consider the process $e^+e^- \rightarrow \Lambda\bar{\Lambda}$. Although the Λ and $\bar{\Lambda}$ hyperons have identical mean lifetimes, their decay times are statistically independent. As a result, the two decay events may be either spacelike or timelike separated, depending on their kinematics. The corresponding spacetime coordinates can be written as (ct_1, vt_1) and $(ct_2, -vt_2)$, respectively, leading to the spacelike-separation condition

$$\Delta s^2 = c^2(t_1 - t_2)^2 - v^2(t_1 + t_2)^2 < 0, \quad (32)$$

as illustrated in Fig. 6. This condition can be rewritten as

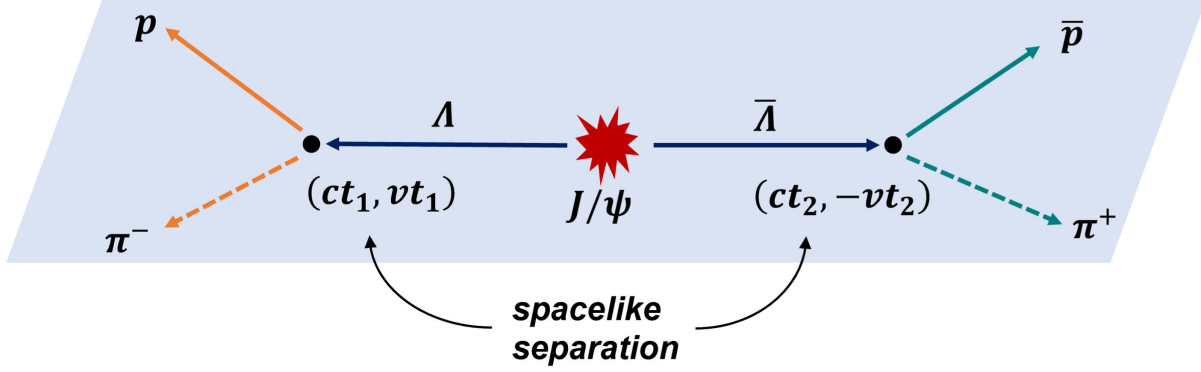


FIG. 6. Schematic diagram of the decay process $J/\psi \rightarrow \Lambda \bar{\Lambda} \rightarrow p \pi^+ \bar{p} \pi^-$. According to the locality loophole, two spacetime coordinates, (ct_1, vt_1) and $(ct_2, -vt_2)$, should be spacelike separation.

$$\frac{|t_1 - t_2|}{t_1 + t_2} < \beta_\Lambda, \quad (33)$$

where $\beta_\Lambda = v/c = \sqrt{1 - 4M^2/s}$ is the hyperon velocity. For the process $J/\psi \rightarrow \Lambda \bar{\Lambda}$, one has $\beta_\Lambda \approx 0.69$, implying that spacelike separation occurs when $|t_1 - t_2|/(t_1 + t_2) < 0.69$.

Since the decay times t_1 and t_2 follow independent exponential distributions [72], the ratio $|t_1 - t_2|/(t_1 + t_2)$ is uniformly distributed in the interval $[0, 1]$, as follows:

$$\frac{|t_1 - t_2|}{t_1 + t_2} \sim \text{Uniform}[0, 1]. \quad (34)$$

As a result, the probability of spacelike separation depends solely on the hyperon velocity,

$$\Pr(\Delta s^2 < 0) = \beta_\Lambda, \quad (35)$$

and can be straightforwardly evaluated for different production energies, as is shown in Fig. 7.

Following previous studies [16,23], both spacelike- and timelike-separated events are incorporated through an effective bound on the Bell inequality. Here, we note that since the spacelike-separation condition in Eq. (33) depends only on the particle velocity, this quantity can

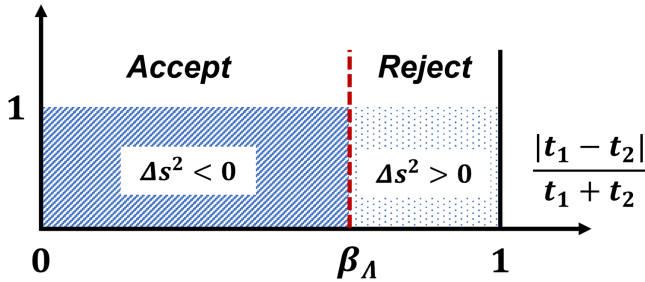


FIG. 7. Fraction of events in which Λ and $\bar{\Lambda}$ decays are spacelike separated is given by β_Λ . The random variable in the lhs of Eq. (33) is a uniform distribution in $[0, 1]$.

be used as a practical criterion in the event reconstruction procedure. In particular, hyperons produced at higher c.m. energies are favored due to their larger velocities; for example, the $\psi(3686)$ resonance is preferable to the J/ψ in this respect. At the same time, the available dataset size must also be considered. For instance, the BESIII J/ψ dataset contains approximately 10×10^9 events, whereas the $\psi(3686)$ dataset comprises about 2.7×10^9 events.

IX. DECOHERENCE AND QUANTUM CORRELATIONS

Quantum decoherence plays a vital role in understanding the quantum nature of physical systems, characterizing their evolution under the influence of environmental interactions [73]. Recent studies have introduced decoherence effects into the description of high-energy processes [74,75]. At collider experiments such as BESIII, the observation of quantum correlations can be challenged by decoherence phenomena [18,76], which arise from interactions between high-energy particles and the detector materials [71,77–80]. In this section, we discuss the role of quantum decoherence in detectors and its potential impact on the measurement of quantum correlations.

We take the decay $\eta_c \rightarrow \Lambda \bar{\Lambda}$ as an example. According to its mean lifetime $\sim 2.6 \times 10^{-10}$ s and $\beta_\Lambda \approx 0.66$, Λ will undergo a secondary decay at an average distance of 5.1 cm from the production vertex. Given that the inner radius of the beam pipe at BESIII is 3.15 cm, it is reasonable to assume that most Λ 's interact with the beam pipe wall and enter the first few layers of the main drift chamber (MDC). When Λ hyperons enter the MDC, their spatial degrees of freedom are modified due to the localization of their trajectories, and their spin orientations can be influenced as well.

Owing to the exponential distribution of their lifetimes, a fraction of short-lived Λ hyperons will decay inside the beam pipe. We exclude these hyperons from our model, as they never interact with any detector materials. In contrast, the longer-lived Λ hyperons traverse through the beam pipe

or other detector components and interact with detector materials. The latter Λ s are the focus of our analysis, and we will describe their spin evolution within the framework of open quantum systems.

In quantum information theory, the evolution of an open quantum system can be described by the Lindblad (or Gorini-Kossakowski-Sudarshan-Lindblad) master equation [81–83],

$$\frac{d\rho}{dt} = -i[\mathbf{H}, \rho] + \sum_{\ell} \left(\mathbf{L}_{\ell} \rho \mathbf{L}_{\ell}^{\dagger} - \frac{1}{2} \{ \mathbf{L}_{\ell}^{\dagger} \mathbf{L}_{\ell}, \rho \} \right), \quad (36)$$

where ρ is the spin density matrix of a particle. Equation (36) contains both the unitary dynamics generated by the effective Hamiltonian $\mathbf{H}$ and the nonunitary dynamics induced by system-environment interactions, represented by the jump operators $\mathbf{L}_{\ell}$.

In principle, the explicit form of the jump operators $\mathbf{L}_{\ell}$ should be derived from a microscopic model of hyperon-material interactions. However, for the purposes of this work, we take a phenomenological approach. Specifically, we model the decoherence of the Λ hyperon via the phase damping channel,

$$\frac{d\rho_{\Lambda}}{dt} = -\frac{\gamma}{2} [\sigma_z, [\sigma_z, \rho_{\Lambda}]], \quad (37)$$

where we omit the unitary term $-i[\mathbf{H}, \rho_{\Lambda}]$ since it does not affect quantum correlations, and take $\sqrt{\gamma}\sigma_z$ as the jump operator.

Here, the z axis is defined along the momentum (flight) direction of the Λ hyperon. As the Λ traverses through the beam pipe or detector materials, it may undergo weak, stochastic interactions with the surrounding medium. These processes typically act as repeated partial measurements of their spin projection along the momentum direction. From the perspective of open quantum systems, such interactions preferentially select the eigenbasis of σ_z as the pointer basis—the basis that remains stable under environmental monitoring. This leads to progressive suppression of coherence between spin-up and spin-down states along the flight direction, while populations in that basis remaining unaffected. Mathematically, this behavior is captured by Eq. (37), which describes a continuous weak measurement of the spin- z component, consistent with the expected symmetry in the rest frame of the hyperon [84].

The phase damping channel in Eq. (37) can be equivalently written in the operator-sum (Kraus) representation:

$$\rho_{\Lambda} \mapsto \mathcal{E}(\rho_{\Lambda}) \equiv \mathbf{K}_0 \rho_{\Lambda} \mathbf{K}_0 + \mathbf{K}_1 \rho_{\Lambda} \mathbf{K}_1, \quad (38)$$

where the Kraus operators are

$$\mathbf{K}_0 = \begin{bmatrix} 1 & 0 \\ 0 & \sqrt{1-\zeta} \end{bmatrix}, \quad \mathbf{K}_1 = \begin{bmatrix} 0 & 0 \\ 0 & \sqrt{\zeta} \end{bmatrix}, \quad (39)$$

with $\sqrt{1-\zeta} \equiv e^{-2\gamma t}$. Acting on the Bloch vector $\mathbf{B}^+ = (B_x^+, B_y^+, B_z^+)$, the channel yields

$$(B_x^+, B_y^+, B_z^+) \mapsto (e^{-2\gamma t} B_x^+, e^{-2\gamma t} B_y^+, B_z^+), \quad (40)$$

which is the solution of the master equation Eq. (37): transverse components decay with the rate 2γ while the longitudinal component is conserved. In NMR terminology, this corresponds to T_2 -type relaxation with $1/T_2 = 2\gamma$. Phase damping is commonly used to model the loss of coherence (information) induced by random scattering processes, analogous to the interactions experienced by a Λ hyperon traversing through detector materials [85].

For a two-qubit system, e.g., $\Lambda\bar{\Lambda}$, we assume that the antiparticle undergoes the same decoherence process as the particle [Eqs. (37) and (38)]. The complete two-qubit channel is then given by $\rho_{\Lambda\bar{\Lambda}} \mapsto \tilde{\rho}_{\Lambda\bar{\Lambda}} \equiv \mathcal{E} \otimes \mathcal{E}(\rho_{\Lambda\bar{\Lambda}})$. Since decoherence corresponds to quantum information leaking into the environment, it reduces the correlation between Λ and $\bar{\Lambda}$.

We can employ the concurrence as a quantitative measure to illustrate how decoherence degrades the entanglement of the $\Lambda\bar{\Lambda}$ system. The numerical results are presented in Fig. 8, which shows that the concurrence decreases monotonically with the decoherence parameter ζ and eventually vanishes in the strong-decoherence limit, reflecting the complete loss of the entanglement between Λ and $\bar{\Lambda}$.

Additionally, in measurements of the process $e^+e^- \rightarrow \Lambda\bar{\Lambda}$, the angular distributions of the secondary decays are used to extract the parameter $\Delta\Phi$ without accounting for decoherence effects. However, because the current experimental results (e.g., $\Delta\Phi = 0.752$ for Λ) have already included decoherence effects, an underestimation of the entanglement relative to its value at the production vertex should be noticed.

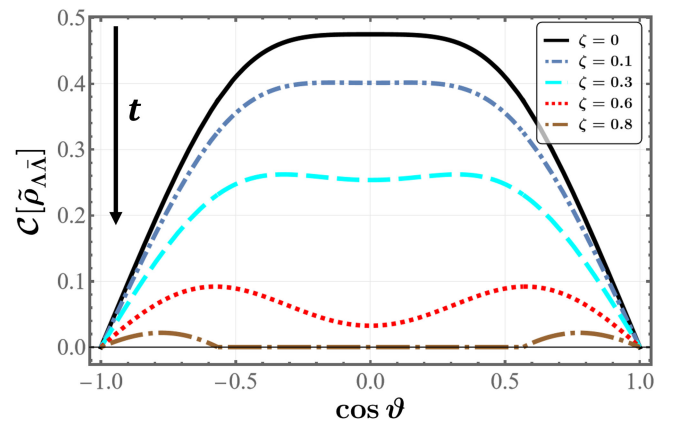


FIG. 8. Entanglement between Λ and $\bar{\Lambda}$ will decrease due to the quantum decoherence effect. The decoherence time t is related to the factor ζ as $\sqrt{1-\zeta} = e^{-2\gamma t}$.

To obtain more accurate values of the quantum correlation and also $\Delta\Phi$, a promising approach is proposed in Ref. [18]: using $\eta_c/\chi_{c0} \rightarrow \Lambda\bar{\Lambda}$ to benchmark the decoherence effect for the Λ hyperon. The above decay channels produce two Bell states from the symmetry:

$$|\Psi^\pm\rangle = \frac{1}{\sqrt{2}}(|\uparrow\downarrow\rangle \pm |\downarrow\uparrow\rangle). \quad (41)$$

By comparing the theoretical predictions in Eq. (41) with experimental results reconstructed via quantum state tomography, one can confirm the presence of decoherence for the Λ hyperon in the detector and identify the specific decoherence channel, such as determining the relaxation time T_2 in a dephasing channel.

Once the quantum channels established from processes like $\eta_c/\chi_{c0} \rightarrow \Lambda\bar{\Lambda}$, the original values of $B_y^\pm$ (without decoherence) can be extracted from experimental measurements after relaxation using Eq. (40). Therefore, the experimental values of $\Delta\Phi$ in $J/\psi \rightarrow \Lambda\bar{\Lambda}$ at BESIII—currently extracted without accounting for decoherence—can be benchmarked. This would make it possible to recover the intrinsic quantum correlations in Λ and $\bar{\Lambda}$.

X. SUMMARY AND OUTLOOK

In this work, we investigated two types of quantum correlations, namely the quantum steering and discord, in hyperon-antihyperon systems in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$, where Y represents octet spin-1/2 hyperons in the ground state: Λ , Σ^+ , Σ^0 , Ξ^- , and Ξ^0 . The analytical expressions were derived for the quantum steering and discord based on the X-shaped density operator for $Y\bar{Y}$ systems, and the results depend on the parameters α_ψ and $\Delta\Phi$ measured in BESIII experiments. We found that the steerability of certain hyperon-antihyperon system is nonzero within a range centered around the scattering angle $\vartheta = 90^\circ$, while the discord is nonzero at all scattering angles. This aligns with the fact that almost all quantum states have nonzero discord [86]. Based on these results, we presented the hierarchy relations among four types of quantum correlations. After that, we examined how electromagnetic form factors influence the correlations in hyperon-antihyperon systems.

The locality loophole and detector-induced decoherence are two major concerns in collider-based tests of quantum correlations. To connect our analysis more closely with realistic collider experiments, we first proposed a practical criterion to identify events with spacelike-separated decays in high-energy collisions. Secondly, we developed a phenomenological model based on the Lindblad master equation to describe quantum decoherence. This framework offers insight into the decoherence mechanism of hyperons and suggests a strategy for extracting more accurate values

of some spin correlation parameters in $J/\psi \rightarrow \Lambda\bar{\Lambda}$ from experimental data.

Recent preliminary measurements indicate an oscillatory behavior in $|G_E/G_M|$ and in $\Delta\Phi$ as functions of the collision energy $\sqrt{s}$, and several theoretical models have been proposed to account for this phenomenon [87–93]. This oscillatory pattern may reflect nontrivial spin correlations in the $Y\bar{Y}$ system. With increasing data statistics and further theoretical development, it may become feasible to systematically investigate the energy dependence of quantum correlations in $Y\bar{Y}$ production across the continuum region of e^+e^- annihilation, beyond the vicinity of specific resonances.

From the perspective of high-energy physics studies, the quantum steering and discord provide new witnesses for quantum correlations in high-energy physics. Quantum properties such as Bell nonlocality, steering, entanglement, and discord may in turn serve as sensitive probes to hadrons' structures such as EMFFs and to fundamental symmetries in their interactions [94]. The distinctive quantum features of hyperon-antihyperon systems also raise the possibility of using quantum correlations to test whether a particle is composite or elementary. Furthermore, studying quantum correlations in hyperon systems provides new insights into the properties of spin, a continuously active field of research [95–99].

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DATA AVAILABILITY

No data were created or analyzed in this study.

APPENDIX A: DETECTING STEERABILITY VIA ENTANGLEMENT

Apart from the steering inequality, there is another method to quantify the steering based on the quantum entanglement [51,52].

Considering a two-qubit state, the steering from Bob to Alice can be witnessed if the density matrix $\rho_{A \leftarrow B}$ defined as

$$\rho_{A \leftarrow B} = \frac{1}{\sqrt{3}}\rho_{AB} + \left(1 - \frac{1}{\sqrt{3}}\right)\rho_A \otimes \frac{1}{2}, \quad (A1)$$

is entangled. If the original ρ_{AB} is a X state, the steered state $\rho_{A \leftarrow B}$ also preserves the X structure

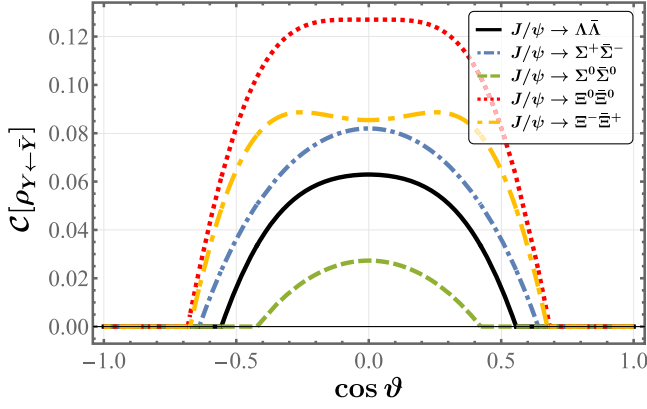


FIG. 9. Results for the concurrence $\mathcal{C}[\rho_{Y\leftarrow\bar{Y}}]$ as functions of $\cos\vartheta$ (ϑ is the scattering angle) in $e^+e^- \rightarrow J/\psi \rightarrow Y\bar{Y}$ with $Y = \Lambda, \Sigma^+, \Sigma^0, \Xi^0$, and Ξ^- in the black solid, blue dash-dotted, green dashed, red dotted, and yellow long dash-dotted lines, respectively.

$$\rho_{A\leftarrow B} = \begin{bmatrix} \frac{1}{\sqrt{3}}\rho_{11} + r & 0 & 0 & \frac{1}{\sqrt{3}}\rho_{14} \\ 0 & \frac{1}{\sqrt{3}}\rho_{22} + r & \frac{1}{\sqrt{3}}\rho_{23} & 0 \\ 0 & \frac{1}{\sqrt{3}}\rho_{23} & \frac{1}{\sqrt{3}}\rho_{33} + s & 0 \\ \frac{1}{\sqrt{3}}\rho_{14} & 0 & 0 & \frac{1}{\sqrt{3}}\rho_{44} + s \end{bmatrix}, \quad (\text{A2})$$

where $r = \frac{3-\sqrt{3}}{6}(\rho_{11} + \rho_{22})$ and $s = \frac{3-\sqrt{3}}{6}(\rho_{33} + \rho_{44})$. Form

the Peres-Horodecki entanglement criterion (also from the Wootters' concurrence), $\rho_{A\leftarrow B}$ is entangled iff the concurrence is greater than zero; i.e., $\mathcal{C}[\rho_{A\leftarrow B}] > 0$. In other words, ρ_{AB} is steerable from Bob to Alice if the concurrence of steered state $\mathcal{C}[\rho_{A\leftarrow B}] > 0$, with

$$\mathcal{C}[\rho_{Y\leftarrow\bar{Y}}] = \frac{1}{2\sqrt{3}} \max \left\{ 0, \begin{aligned} &|t_1 - t_2| - \sqrt{(\sqrt{3} - t_3)^2 - (\sqrt{3} - 1)^2 a^2}, \\ &|t_1 + t_2| - \sqrt{(\sqrt{3} + t_3)^2 - (\sqrt{3} + 1)^2 a^2} \end{aligned} \right\}. \quad (\text{A3})$$

The results for $\mathcal{C}[\rho_{Y\leftarrow\bar{Y}}]$ are shown in Fig. 9. We can see that detecting steerability by entanglement has a similar result compared with CJWR inequality—the $Y\bar{Y}$ state can be steered within the range around the transverse scattering angle. However, some details are different. But the entanglement criterion is also a sufficient but not necessary condition for the quantum steering.

APPENDIX B: QUANTUM CORRELATIONS IN $Y\bar{Y}$ SYSTEMS THROUGH $\psi(3686)$

In this Appendix, we calculate the quantum correlations in $Y\bar{Y}$ systems through $\psi(3686)$ in e^+e^- annihilation, using the experimental data for the parameters α_ψ and $\Delta\Phi$ in Table III. The results are shown in Fig. 10. This section is complementary to Sec. V and aims to extend the investigation from J/ψ to $\psi(3686)$.

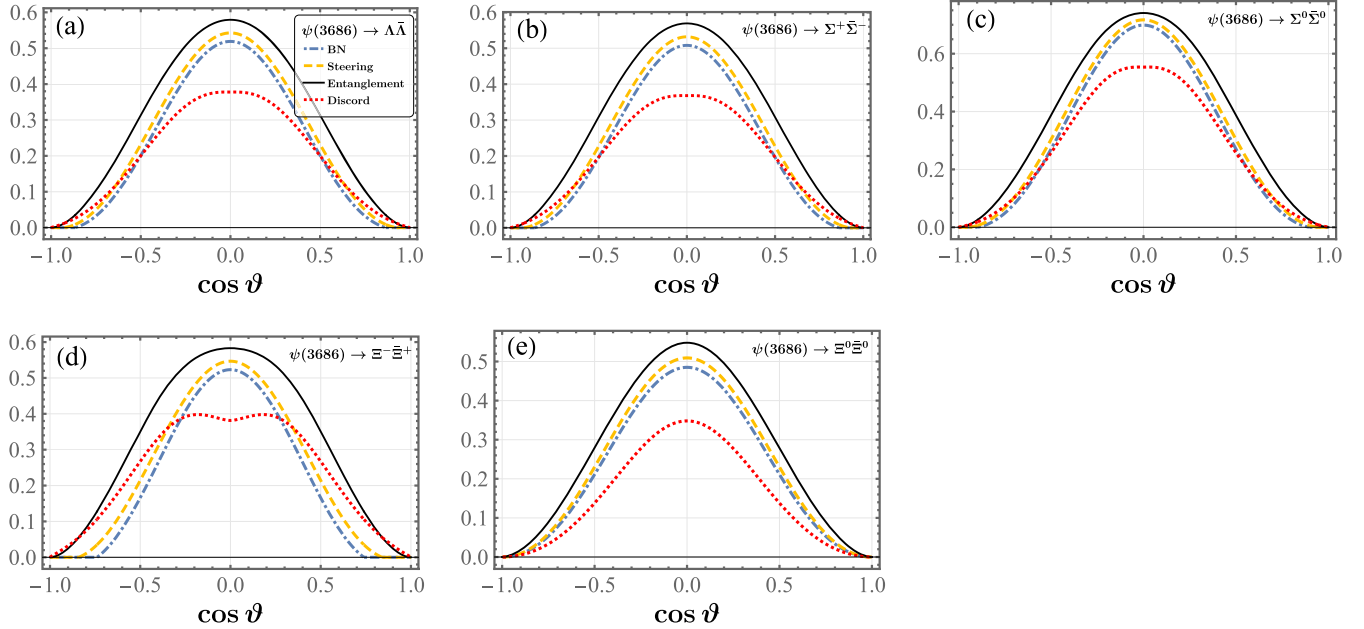


FIG. 10. Four types of quantum correlations as functions $\cos\vartheta$ in $e^+e^- \rightarrow \psi(3686) \rightarrow Y\bar{Y}$. The panels from (a)–(e) correspond to $\Lambda, \Sigma^+, \Sigma^0, \Xi^-,$ and Ξ^0 , respectively. In each panel, the results for the Bell nonlocality $\mathcal{B}$, steering $\mathcal{S}$, entanglement $\mathcal{E}$, and discord $\mathcal{D}$ are shown in blue dot-dashed, yellow dashed, black solid, and red dotted lines, respectively.

TABLE III. Some parameters in $e^+e^- \rightarrow \psi(3686) \rightarrow Y\bar{Y}$, where $Y\bar{Y}$ is a pair of ground-state octet hyperons.

	Branching ratio ($\times 10^{-4}$)	α_ψ	$\Delta\Phi/\text{rad}$	Ref.
$\psi(3686) \rightarrow \Lambda\bar{\Lambda}$	3.81 ± 0.13	$0.69 \pm 0.07 \pm 0.02$	$0.40 \pm 0.14 \pm 0.03$	[46,100]
$\psi(3686) \rightarrow \Sigma^+\bar{\Sigma}^-$	2.82 ± 0.09	$0.682 \pm 0.030 \pm 0.011$	$0.397 \pm 0.07 \pm 0.014$	[43,46]
$\psi(3686) \rightarrow \Sigma^0\bar{\Sigma}^0$	2.35 ± 0.09	$0.814 \pm 0.028 \pm 0.028$	$0.512 \pm 0.085 \pm 0.034$	[44,46]
$\psi(3686) \rightarrow \Xi^-\bar{\Xi}^+$	2.87 ± 0.11	$0.693 \pm 0.048 \pm 0.049$	$0.667 \pm 0.111 \pm 0.058$	[46,101]
$\psi(3686) \rightarrow \Xi^0\bar{\Xi}^0$	2.3 ± 0.4	$0.665 \pm 0.086 \pm 0.081$	$-0.050 \pm 0.150 \pm 0.020$	[46,48]

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