

SMEFT effects on spin correlations and entanglement at NLO QCD in diboson production at hadron colliders

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We perform for the first time a full study of spin correlations in inclusive production of WZ boson pairs at the LHC with leptonic decays in the presence of next-to-leading-order QCD corrections and of effects from a dimension-six operator in the Standard-Model effective field theory (SMEFT) modifying the electroweak triple-gauge coupling. We carry out the complete quantum-state tomography of the diboson system and relate its results to common purity and spin-entanglement markers, highlighting the sizeable impact of both QCD corrections and SMEFT insertions. Additionally, we show how a naïve truncation at dimension six in the SMEFT expansion of the spin-density matrix can lead to a cumbersome spin interpretation of the quantum-tomography results.

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I. INTRODUCTION

The strong interest in measuring observables inspired by quantum information (QI) at the Large Hadron Collider (LHC) [1,2], further increased by the entanglement studies by ATLAS and CMS in the top-quark sector [3–5], is not limited to the case of two-level spin states (qubits) but spans also higher dimensionalities and in particular the case of three-level spin states (qutrits). Recently, a large number of phenomenological studies have been devoted to Higgs-boson decays [6–20] and to diboson production at colliders [21–35]. However, in spite of the availability of simulation tools at next-to-leading-order (NLO) accuracy, just a small share of works have highlighted the importance of the inclusion of higher-order corrections in the strong (QCD) and electroweak (EW) couplings [17–19,27,30,32] for an accurate evaluation of spin-correlation coefficients and QI-inspired observables that depend on them. Furthermore, the investigation of entanglement and Bell nonlocality markers in the presence of anomalous couplings has been mostly limited to boson pairs produced in a scalar-boson decay [9,10,13,16,17]. An analysis of Standard-Model effective field theory (SMEFT) effects on the entanglement markers for inclusive diboson production at colliders has been carried in a simplified framework

(no decay simulation, leading-order only) in Ref. [24]. Amongst diboson production channels, ZZ has attracted the greatest attention within the LHC community, owing to the high signal purity in the four-lepton decay channel, while only partial results for the quantum tomography of inclusive WZ production at NLO QCD were obtained in the Standard Model (SM) [27] [see Section (4.1) therein] and in the SMEFT [36] [see Section (4.7) therein]. A broad and realistic study of the WZ spin-density matrix and related QI-inspired observables is still missing in the literature. As a further motivation, the spin structure of WZ production has been studied to high perturbative accuracy for what concerns single-boson and joint polarization fractions in the so-called polarized-template method [37–39]. On the contrary, the calculation of off-diagonal entries of the spin-density matrix for WZ pairs is typically limited to the lowest order.

In this work we achieve for the first time a complete analysis of spin correlations for WZ inclusive production at the LHC in the presence of NLO QCD corrections and of effects from a SMEFT dimension-six operator modifying the triple-gauge coupling. On top of the full quantum tomography, we discuss the effects of higher-order QCD corrections and SMEFT effects on relevant markers for the purity and the entanglement associated to the spin structure of the WZ system.

The article is structured as follows. In Sec. II we introduce the general aspects of single- and two-boson spin-density matrices and their relation to the extraction of spin correlations via quantum tomography. In Sec. III we highlight potential issues in the interpretation of spin correlations that arise in the presence of new-physics effects described by dimension-six SMEFT operators. The numerical results are discussed in Sec. IV, while in Sec. V we draw our conclusions.

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II. GENERALITIES

At tree level but also including higher orders in QCD, if one assumes a two-body decay of an intermediate weak boson into leptons, the differential distribution in the decay angles θ, ϕ has a compact analytic form [40,41], which can be written as an expansion in the orthonormal basis of real spherical harmonics $Y_{lm}(\Omega)$ up to rank $l = 2$ [7],

$$\frac{d\sigma}{d\Omega dX} \Big/ \frac{d\sigma}{dX} = \frac{1}{4\pi} + \sum_{l=1}^2 \sum_{m=-l}^l \alpha_{lm}(X) Y_{lm}(\Omega),$$

$$d\Omega = d\phi d\cos\theta, \quad (1)$$

where X is a generic observable that is independent of the decay angular measure, e.g., the transverse momentum of the decaying boson. Since we are interested in diboson systems, we define decay angles in the so-called modified helicity coordinate system [42], i.e., θ, ϕ are the polar and azimuthal angles of the positively charged lepton in the

corresponding boson rest frame, computed taking as reference axis for the polar angle θ the boson spatial direction in the diboson center-of-mass (CM) frame. In the absence of kinematic constraints on the decay products, the α_{lm} coefficients can be extracted by projecting onto corresponding spherical harmonics,

$$\alpha_{lm}(X) = \frac{\int d\Omega \frac{d\sigma}{d\Omega dX} Y_{lm}(\Omega)}{\int d\Omega \frac{d\sigma}{d\Omega dX}}. \quad (2)$$

To relate the angular coefficients to the spin-density matrix for a single boson V we rewrite Eq. (1) as,

$$\frac{d\sigma}{d\Omega dX} \Big/ \frac{d\sigma}{dX} = \frac{3}{4\pi} \sum_{i,j=1}^3 \rho_{V,ij}(X) \Gamma_{ij}(\Omega), \quad (3)$$

where Γ is the decay matrix for the two-body decay to massless leptons [21,43],

$$\Gamma(\Omega) = \frac{1}{4} \begin{pmatrix} 1 + 2\eta_\ell \cos\theta + \cos^2\theta & \sqrt{2} \sin\theta(\eta_\ell + \cos\theta)e^{i\phi} & (1 - \cos^2\theta)e^{2i\phi} \\ \sqrt{2} \sin\theta(\eta_\ell + \cos\theta)e^{-i\phi} & 2(1 - \cos^2\theta) & \sqrt{2} \sin\theta(\eta_\ell - \cos\theta)e^{i\phi} \\ (1 - \cos^2\theta)e^{-2i\phi} & \sqrt{2} \sin\theta(\eta_\ell - \cos\theta)e^{-i\phi} & 1 - 2\eta_\ell \cos\theta + \cos^2\theta \end{pmatrix}, \quad (4)$$

and ρ_V is a linear combination of the identity $\mathbb{I}_3$ and of the irreducible tensor representations for qutrits T_{lm} [7,17],¹

$$\rho_V(X) = \frac{1}{3} \left[\mathbb{I}_3 + \sum_{l=1}^2 \kappa_l \sum_{m=-l}^l \alpha_{lm}(X) T_{lm} \right], \quad \kappa_1 = \frac{\sqrt{8\pi}}{\eta_\ell}, \quad \kappa_2 = \sqrt{40\pi}. \quad (5)$$

We note that the potential dependence on the kinematics of the decaying boson (X) enters the production spin-density matrix ρ_V but not the decay matrix Γ . However, it was recently shown [32] that a mild dependence on the boson transverse momentum can be present in the spin-analyzing power η_ℓ in Eq. (4), owing to lepton dressing starting at NLO EW. Since in this work we only consider QCD corrections we can safely assume that Γ is completely independent of any X other than the decay angles. Furthermore, the analytic structure of Eq. (4) is valid in the presence of higher-order QCD corrections, and to large extent also of radiative EW corrections, up to the modification of the numerical value of the spin-analyzing power η_ℓ [32].

The spin-density matrix in Eq. (5) (but also its generalization to a system of two or more bosons) is Hermitian, has unit trace, and is semipositive definite. Being in general a mixture of pure-state projectors, it can be written in terms of helicity amplitudes [1] as follows,

$$\rho_{V,ij} = \frac{\sum_x \mathcal{A}_{i,x} \mathcal{A}_{j,x}^*}{\sum_x \sum_{a=1}^3 |\mathcal{A}_{a,x}|^2}, \quad \text{Tr}[\rho_V] = 1, \quad \frac{1}{3} \leq \text{Tr}[\rho_V^2] \leq 1, \quad (6)$$

¹The T_{lm} matrices are defined such that $T_{l,-m} = (-1)^m T_{lm}^\dagger$, and read,

$$T_{11} = \sqrt{\frac{3}{2}} \begin{pmatrix} 0 & -1 & 0 \\ 0 & 0 & -1 \\ 0 & 0 & 0 \end{pmatrix}, \quad T_{10} = \sqrt{\frac{3}{2}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -1 \end{pmatrix}, \quad T_{22} = \sqrt{3} \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix},$$

$$T_{21} = \sqrt{\frac{3}{2}} \begin{pmatrix} 0 & -1 & 0 \\ 0 & 0 & 1 \\ 0 & 0 & 0 \end{pmatrix}, \quad T_{20} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & -2 & 0 \\ 0 & 0 & 1 \end{pmatrix}.$$

where the sum over x understands sums and averages over all degrees of freedom other than the helicity of the considered boson. We note that $\text{Tr}[\rho_V^2] = 1$ is achieved for pure states, while $\text{Tr}[\rho_V^2] \geq \frac{1}{3}$ is a direct consequence of the Cauchy-Schwartz inequality. From Eq. (5), we obtain the following expression for the trace of ρ_V^2 in terms of α_{lm} coefficients (making the dependence on independent variables X implicit),

$$\text{Tr}[\rho_V^2] = \frac{1}{3} \left[1 + \sum_{l=1}^2 \sum_{m=-l}^l (\kappa_l \alpha_{lm})^2 \right], \quad (7)$$

which can be used to assess the purity of generic mixed states.

The generalization of Eq. (1) to the case of two EW bosons reads [7,17,27],

$$\begin{aligned} \frac{d\sigma}{d\Omega_1 d\Omega_2 dX} / \frac{d\sigma}{dX} = & \frac{1}{(4\pi)^2} + \frac{1}{4\pi} \sum_{l=1}^2 \sum_{m=-l}^l \alpha_{lm}^{(1)}(X) Y_{lm}(\Omega_1) + \frac{1}{4\pi} \sum_{l'=1}^2 \sum_{m'=-l'}^{l'} \alpha_{l'm'}^{(2)}(X) Y_{l'm'}(\Omega_2) \\ & + \sum_{l,l'=1}^2 \sum_{m=-l}^l \sum_{m'=-l'}^{l'} \gamma_{lm'l'm'}(X) Y_{lm}(\Omega_1) Y_{l'm'}(\Omega_2), \end{aligned} \quad (8)$$

where Ω_1, Ω_2 represent the decay angles associated to the first and second boson, respectively. All polarization ($\alpha_{lm}^{(1)}, \alpha_{lm}^{(2)}$) and spin-correlation ($\gamma_{lm'l'm'}$) coefficients can be extracted from Eq. (8) by projecting onto one or two spherical harmonics, in the same fashion as in Eq. (2). Following the same reasoning as above, the four-dimensional angular distribution in Eq. (8) can also be written as a combination of the spin-density matrix ρ for two qutrits and two decay matrices, extending Eq. (3),

$$\frac{d\sigma}{d\Omega_1 d\Omega_2 dX} / \frac{d\sigma}{dX} = \left(\frac{3}{4\pi} \right)^2 \sum_{i,i',j,j'=1}^3 \rho_{ii'jj'}(X) \Gamma_{ii'}(\Omega_1) \Gamma_{jj'}(\Omega_2), \quad (9)$$

where the decay matrices have the form of Eq. (4) and ρ is now a Hermitian, unit-trace, 9×9 matrix that reads,

$$\begin{aligned} \rho(X) = & \frac{1}{9} \left[\mathbb{I}_3 \otimes \mathbb{I}_3 + \sum_{l=1}^2 \kappa_l^{(1)} \sum_{m=-l}^l \alpha_{lm}^{(1)}(X) T_{lm} \otimes \mathbb{I}_3 + \sum_{l'=1}^2 \kappa_{l'}^{(2)} \sum_{m'=-l'}^{l'} \alpha_{l'm'}^{(2)}(X) \mathbb{I}_3 \otimes T_{l'm'} \right. \\ & \left. + \sum_{l,l'=1}^2 \kappa_l^{(1)} \kappa_{l'}^{(2)} \sum_{m=-l}^l \sum_{m'=-l'}^{l'} \gamma_{lm'l'm'}(X) T_{lm} \otimes T_{l'm'} \right], \\ \kappa_1^{(i)} = & \frac{\sqrt{8\pi}}{\eta_{\ell,i}}, \quad \kappa_2^{(i)} = \sqrt{40\pi}. \end{aligned} \quad (10)$$

Note that for $V_1 = W, V_2 = Z$ the two spin-analyzing powers are different, $\eta_{\ell,W} \neq \eta_{\ell,Z}$. Although we do not write it explicitly, the ρ matrix for two bosons can be easily expressed in terms of helicity amplitudes similarly to Eq. (6) by making two helicity indices explicit (one for each boson). The purity of the diboson state can be assessed by taking the trace of the square of the ρ matrix, which reads [17],

$$\text{Tr}[\rho^2] = \frac{1}{9} \left[1 + \sum_{l=1}^2 \sum_{m=-l}^l (\kappa_l^{(1)} \alpha_{lm}^{(1)})^2 + \sum_{l'=1}^2 \sum_{m'=-l'}^{l'} (\kappa_{l'}^{(2)} \alpha_{l'm'}^{(2)})^2 + \sum_{l,l'=1}^2 \sum_{m=-l}^l \sum_{m'=-l'}^{l'} (\kappa_l^{(1)} \kappa_{l'}^{(2)} \gamma_{lm'l'm'})^2 \right]. \quad (11)$$

From Eqs. (5) and (10) it is clear that by extracting the coefficients with $\ell \leq 2$ from angular distributions in Eqs. (1) and (8), one can entirely reconstruct the structure of the spin-density matrix for one (ρ_V) and two (ρ) EW bosons, or qutrits. This approach goes by the name of quantum-state tomography. From the full knowledge of the spin-density matrix one can construct suitable markers for entanglement, Bell nonlocality, and other QI-inspired

observables. In this work we focus on spin entanglement. The level of entanglement for two-qutrit systems can be evaluated through the concurrence [44–47] for which a closed analytic expression only exists in special configurations. In the case of generic mixed states, the concurrence $\mathcal{C}$ is defined through an optimization procedure and only lower and upper bounds are known [45,46]. These can be expressed as [17],

$$2 \max \left(\text{Tr}[\rho^2] - \text{Tr}[\rho_{V_1}^2], \text{Tr}[\rho^2] - \text{Tr}[\rho_{V_2}^2] \right) \leq C^2$$

$$\leq 2 \min \left(1 - \text{Tr}[\rho_{V_1}^2], 1 - \text{Tr}[\rho_{V_2}^2] \right), \quad (12)$$

where the trace of squared spin-density matrices for individual bosons and for the boson pair take the form of Eqs. (7) and (11), respectively. If the lower bound in Eq. (12) is larger than zero, then the two-boson spin state is entangled.

III. SMEFT EFFECTS

The whole treatment above applies as long as the probabilistic interpretation of the spin-density matrix (semipositivity) is not jeopardized. This holds independently of the specific dynamical theory underlying the EW-boson production, whether the SM or any beyond-the-SM theory.

We now focus on potential new-physics effects modeled in the framework of the SMEFT. For simplicity, but without loss of generality of our conclusions, we consider one of the least constrained CP -even operators leading to anomalous triple-gauge couplings (aTGC) in the EW sector and often studied in the context of diboson processes [36,48–63],

$$\mathcal{L}_{\text{SMEFT}} = \mathcal{L}_{\text{SM}} + \frac{C_W}{\Lambda^2} \mathcal{Q}_W,$$

$$\mathcal{Q}_W = \epsilon_{ijk} W_\mu^{i,\nu} W_\nu^{j,\lambda} W_\lambda^{k,\mu}, \quad (13)$$

where C_W is the Wilson coefficient (WC) associated to the $\mathcal{Q}_W$ operator and Λ is the scale of new physics. Up to Λ^{-4} , and with a single insertion of the considered operator, a standard observable, e.g., the transverse momentum of the EW boson, would receive contributions from squared SM (Λ^0) amplitudes, from the linear interference between SM and EFT amplitudes (Λ^{-2}), and from squared EFT amplitudes (Λ^{-4}),

$$\frac{d\sigma}{dX} = \frac{d\sigma^{(4)}}{dX} + \xi \frac{d\sigma^{(6)}}{dX} + \xi^2 \frac{d\sigma^{(8)}}{dX}, \quad (14)$$

where we define $\xi = C_W/\Lambda^2$ and the canonical dimensions of the three contributions are used in the labels [(4) for the SM, (6) for the linear SM–EFT interference, and (8) for the quadratic EFT term]. The impact on the angular coefficients is however more intricate, as the dependence on the SMEFT effects enters both in the numerator and in the denominator of Eq. (2) and therefore they are extracted as [62,64],

$$\alpha_{lm}(X) = \frac{\int d\Omega \left[\frac{d\sigma^{(4)}}{d\Omega dX} + \xi \frac{d\sigma^{(6)}}{d\Omega dX} + \xi^2 \frac{d\sigma^{(8)}}{d\Omega dX} \right] Y_{lm}(\Omega)}{\int d\Omega \left[\frac{d\sigma^{(4)}}{d\Omega dX} + \xi \frac{d\sigma^{(6)}}{d\Omega dX} + \xi^2 \frac{d\sigma^{(8)}}{d\Omega dX} \right]}. \quad (15)$$

In the SMEFT the generic amplitude $\mathcal{A}_i$ reads $\mathcal{A}_i^{(4)} + \xi \mathcal{A}_i^{(6)}$, therefore the spin-density matrix in Eq. (6) becomes,

$$\rho_{V,ij}^{(8)} = \frac{\sum_x (\mathcal{A}_{i,x}^{(4)} + \xi \mathcal{A}_{i,x}^{(6)}) (\mathcal{A}_{j,x}^{(4)*} + \xi \mathcal{A}_{j,x}^{(6)*})}{\sum_x \sum_{a=1}^3 |\mathcal{A}_{a,x}^{(4)} + \xi \mathcal{A}_{a,x}^{(6)}|^2} = \frac{\sum_x \left[(\mathcal{A}_{i,x}^{(4)} \mathcal{A}_{j,x}^{(4)*}) + \xi (\mathcal{A}_{i,x}^{(4)} \mathcal{A}_{j,x}^{(6)*} + \mathcal{A}_{i,x}^{(6)} \mathcal{A}_{j,x}^{(4)*}) + \xi^2 (\mathcal{A}_{i,x}^{(6)} \mathcal{A}_{j,x}^{(6)*}) \right]}{\sum_x \sum_{a=1}^3 \left[|\mathcal{A}_{a,x}^{(4)}|^2 + 2\xi \Re(\mathcal{A}_{a,x}^{(4)} \mathcal{A}_{a,x}^{(6)*}) + \xi^2 |\mathcal{A}_{a,x}^{(6)}|^2 \right]}, \quad (16)$$

where all terms in the SMEFT ξ expansion are retained, including quadratic ones. Just like Eq. (6), the SMEFT spin-density matrix in Eq. (16) is Hermitian, has unit trace, and is semipositive definite.

The SMEFT expansion is often studied up to linear order in ξ , i.e., excluding any effect that is formally beyond order Λ^{-2} . This requires one to ignore quadratic terms from Eq. (15), leading to,

$$\alpha_{lm}(X) = \frac{\int d\Omega \left[\frac{d\sigma^{(4)}}{d\Omega dX} + \xi \frac{d\sigma^{(6)}}{d\Omega dX} \right] Y_{lm}(\Omega)}{\int d\Omega \left[\frac{d\sigma^{(4)}}{d\Omega dX} + \xi \frac{d\sigma^{(6)}}{d\Omega dX} \right]}. \quad (17)$$

If the same is done for the spin-density matrix, the quadratic contributions in ξ have to be discarded both in the numerator and in the denominator of Eq. (16), leading to,

$$\rho_{V,ij}^{(6)} = \frac{\sum_x \left[(\mathcal{A}_{i,x}^{(4)} \mathcal{A}_{j,x}^{(4)*}) + \xi (\mathcal{A}_{i,x}^{(4)} \mathcal{A}_{j,x}^{(6)*} + \mathcal{A}_{i,x}^{(6)} \mathcal{A}_{j,x}^{(4)*}) \right]}{\sum_x \sum_{a=1}^3 \left[|\mathcal{A}_{a,x}^{(4)}|^2 + 2\xi \Re(\mathcal{A}_{a,x}^{(4)} \mathcal{A}_{a,x}^{(6)*}) \right]}. \quad (18)$$

While the ρ_V matrix in Eq. (18) is Hermitian and has unit trace, it is not anymore semipositive definite, leading to important consequences in the interpretation of the angular analysis.

To show that the spin-density matrix at dimension-six in SMEFT like in Eq. (18) is ill defined, we tailor the discussion to the simplified case of a pure state, for which there is no sum over additional degrees of freedom x . It is straightforward to show that the $\rho_V^{(8)}$ matrix for a pure state,

$$\rho_{V,ij}^{(8,\text{pure})} = \frac{(\mathcal{A}_i^{(4)} + \xi \mathcal{A}_i^{(6)})(\mathcal{A}_j^{(4)*} + \xi \mathcal{A}_j^{(6)*})}{\sum_{a=1}^3 |\mathcal{A}_a^{(4)} + \xi \mathcal{A}_a^{(6)}|^2} = \frac{(\mathcal{A}_i^{(4)} \mathcal{A}_j^{(4)*}) + \xi(\mathcal{A}_i^{(4)} \mathcal{A}_j^{(6)*} + \mathcal{A}_i^{(6)} \mathcal{A}_j^{(4)*}) + \xi^2(\mathcal{A}_i^{(6)} \mathcal{A}_j^{(6)*})}{\sum_{a=1}^3 (|\mathcal{A}_a^{(4)}|^2 + 2\xi \Re(\mathcal{A}_a^{(4)} \mathcal{A}_a^{(6)*}) + \xi^2 |\mathcal{A}_a^{(6)}|^2)}, \quad (19)$$

is a projector ($\text{Tr}[\rho_V^2] = 1$) for any value of ξ . If we take Eq. (18) and tailor it to a pure state, we get,

$$\rho_{V,ij}^{(6,\text{pure})} = \frac{(\mathcal{A}_i^{(4)} \mathcal{A}_j^{(4)*}) + \xi(\mathcal{A}_i^{(4)} \mathcal{A}_j^{(6)*} + \mathcal{A}_i^{(6)} \mathcal{A}_j^{(4)*})}{\sum_{a=1}^3 (|\mathcal{A}_a^{(4)}|^2 + 2\xi \Re(\mathcal{A}_a^{(4)} \mathcal{A}_a^{(6)*}))}, \quad (20)$$

then the trace of its square reads,

$$\begin{aligned} \text{Tr}[\rho_V^{(6,\text{pure})2}] &= \frac{\sum_{i,j=1}^3 |(\mathcal{A}_i \mathcal{A}_j^*) + (\mathcal{A}_i \mathcal{F}_j^* + \mathcal{F}_i \mathcal{A}_j^*)|^2}{\left[\sum_{a=1}^3 (|\mathcal{A}_a|^2 + (\mathcal{A}_a \mathcal{F}_a^* + \mathcal{A}_a^* \mathcal{F}_a)) \right]^2} \\ &= \frac{\sum_{i,j=1}^3 \left[|\mathcal{A}_i|^2 |\mathcal{A}_j|^2 + 2|\mathcal{A}_i|^2 (\mathcal{A}_j \mathcal{F}_j^* + \mathcal{A}_j^* \mathcal{F}_j) + 2|\mathcal{A}_i|^2 |\mathcal{F}_j|^2 + \mathcal{A}_i \mathcal{A}_j \mathcal{F}_i^* \mathcal{F}_j^* + \mathcal{A}_i^* \mathcal{A}_j^* \mathcal{F}_i \mathcal{F}_j \right]}{\sum_{i,j=1}^3 \left[|\mathcal{A}_i|^2 |\mathcal{A}_j|^2 + 2|\mathcal{A}_i|^2 (\mathcal{A}_j \mathcal{F}_j^* + \mathcal{A}_j^* \mathcal{F}_j) + 2\mathcal{A}_i \mathcal{A}_j^* \mathcal{F}_i^* \mathcal{F}_j + \mathcal{A}_i \mathcal{A}_j \mathcal{F}_i^* \mathcal{F}_j^* + \mathcal{A}_i^* \mathcal{A}_j^* \mathcal{F}_i \mathcal{F}_j \right]} \end{aligned} \quad (21)$$

where we have renamed $\mathcal{A}^{(4)} \rightarrow \mathcal{A}$ and $\xi \mathcal{A}^{(6)} \rightarrow \mathcal{F}$. Owing to the Cauchy-Schwartz inequality,

$$\left(\sum_{i=1}^3 |\mathcal{A}_i|^2 \right) \left(\sum_{i=1}^3 |\mathcal{F}_i|^2 \right) \geq \left| \sum_{i=1}^3 \mathcal{A}_i \mathcal{F}_i^* \right|^2, \quad (22)$$

in Eq. (21) the numerator can exceed or equal the denominator, therefore,

$$\text{Tr}[\rho_V^{(6,\text{pure})2}] \geq 1. \quad (23)$$

This implies that, independently of the helicity rules in the SM and in the EFT amplitudes, the truncation of the density-matrix SMEFT expansion at linear level does not allow for a straightforward interpretation as a suitable spin-density matrix. Since such an ill definition of the spin-density matrix holds for a pure state, it cannot be excluded that the same holds for mixed states as well [see Eq. (18)], especially in kinematic regions and dynamical regimes where the purity of the state is high, as we will see in the next Section. On the contrary, the spin-density definitions in Eqs. (16) and (19) correctly satisfy

$$\text{Tr}[\rho_{V,ij}^{(8)}] \leq 1, \quad \text{Tr}[\rho_{V,ij}^{(8,\text{pure})2}] = 1, \quad (24)$$

in arbitrary kinematic and dynamical regimes. As a natural consequence, also the spin-density matrix for a diboson (two-qudit) state is affected by similar issues, when truncating the SMEFT expansion at linear level in the WC of

SMEFT operators. In such case, as for a single boson [see Eq. (18)], the interpretation of the angular coefficients as spin correlations of the diboson system becomes cumbersome. In particular, we will see from the numerical results that the trace in Eq. (11) can exceed the value of one for certain kinematical regimes and certain values of the WC under consideration. We note that similar conclusions were drawn for dipole operators in $\tau^+ \tau^-$ LHC production [65] using the high-energy growth of operators leading to SMEFT interference resurrection (due to different helicity rules) to motivate the violation of perturbative unitarity in the ρ matrix at energy scales lower than Λ . At variance with Ref. [65], our derivation, although applied to a pure-gauge operator, which also leads to interference resurrection in the bosonic sector, is more general and uniquely relying on general aspects of the ρ matrix (Hermiticity, unit trace, semipositive definiteness). A discussion of the effect of the EFT truncation on the ρ -matrix positivity in the case of four-fermion operators affecting $e^+ e^- \rightarrow t \bar{t}$ can also be found in Ref. [66], which appeared concurrently with our work.

IV. NUMERICAL RESULTS

We consider the inclusive production of WZ pairs in Run-3 LHC proton-proton collisions ($\sqrt{s} = 13.6$ TeV). Although it was shown [27,30,32] that the full off-shell picture of angular coefficients in diboson processes is correctly reproduced by factorized calculations with intermediate on-shell bosons, we perform all calculations in the

double-pole approximation with the newly released DI-BOSON package [36,67,68] in the Powheg-Box framework [69–72], ensuring that intermediate spin-1 bosons are on shell. Specifically we consider the decay channel with three massless charged leptons and one neutrino,

$$pp \rightarrow W^+(\rightarrow e^+\nu_e)Z(\rightarrow \mu^+\mu^-). \quad (25)$$

The calculation is carried at both LO and at NLO QCD perturbative accuracy, while parton-shower effects are not considered, since it was proved that they have a mild impact on the angular coefficients [36]. We use the default EW input parameters [36] that are compliant with the recent PDG review [73]. The WC for the Q_W operator [see Eq. (13)] is initially set to $C_W/\Lambda^2 = 1 \text{ TeV}^{-2}$ and then varied between +1 and −1. The tree-level and one-loop SM and SMEFT amplitudes are computed with the `Recola 2` [74] library. The subtraction of QCD infrared singularities is achieved in the Frixione-Kunst-Signer scheme [75]. Both the factorization and the renormalization scales are set to the arithmetic average of the W- and Z-boson pole masses. The parton luminosities in the proton parton-distribution function (PDFs), as well as the QCD coupling α_s , are evaluated through the LHAPDF interface [76]. The NNPDF31_NLO_AS_0118 PDF set [77] is used.

For our phenomenological study we consider an inclusive setup,

$$81 \text{ GeV} < M_{\mu^+\mu^-} < 101 \text{ GeV}, \quad (26)$$

and a boosted one,

$$\begin{aligned} 81 \text{ GeV} < M_{\mu^+\mu^-} < 101 \text{ GeV}, \\ p_{T,Z} > 200 \text{ GeV}, \quad p_{T,WZ} < 70 \text{ GeV}. \end{aligned} \quad (27)$$

The two are simplified versions of ATLAS setups for polarization studies [78,79], obtained by retaining only the selections affecting the boson kinematics and not the one of individual decay products. This allows one to pursue the extraction of polarization and spin-correlation coefficients in the same fashion as in Eqs. (2) and (15). In the realistic case of an extraction from LHC data, an extrapolation of angular distributions from fiducial to uncut setups is required to undo the selections that are necessarily applied to decay products in any LHC event selection. This extrapolation is different for SM and beyond-the-SM underlying dynamics, owing to the model dependence of cut effects on polar- and azimuthal-angle distributions [27].

In the following we dub SMEFT6 the predictions that include EFT effects up to dimension-six interference, i.e., single-boson and diboson angular coefficients are extracted as in Eq. (17). We dub SMEFT8 all predictions including both interference and quadratic terms in the EFT expansion, with the coefficients extracted as in Eq. (15).

We first present the $\alpha_{lm}^{(1)}$ angular coefficients, which fully determine the spin-density matrix for the W boson. The numerical results for coefficients that are numerically not compatible with zero are shown in Table I for both considered setups. The NLO QCD corrections mostly affect the $l = 2$ coefficients, while the changes in $\alpha_{10}^{(1)}$ are more moderate. This can be qualitatively explained observing that (see e.g. Refs. [37,62]) at NLO the large gluon-quark contribution changes the LO spin balance in the quark-antiquark channel, leading to a slight enhancement in f_L (f_L being the longitudinal polarization fraction associated to the W) while keeping the left-right balance similar to the one at LO. Since the α_{20} coefficient is directly proportional to $(1 - 3f_L)$, while α_{10} is proportional to the difference between the left and right modes ($f_- - f_+$), it

TABLE I. LO and NLO QCD polarization coefficients for the W boson in the inclusive and boosted setups described in Eqs. (26) and (27), respectively. SM, SMEFT6, and SMEFT8 results were obtained applying Eqs. (2), (17), (15), respectively, and using $C_W/\Lambda^2 = 1 \text{ TeV}^{-2}$. Uncertainties in parentheses are statistical errors from numerical integration.

Coeff.	LO	NLO	LO	NLO	LO	NLO
	SM		SMEFT6		SMEFT8	
Inclusive						
$\kappa_1^{(1)} \cdot \alpha_{10}^{(1)}$	−0.225(1)	−0.234(1)	−0.222(1)	−0.234(1)	−0.185(1)	−0.207(1)
$\kappa_2^{(1)} \cdot \alpha_{20}^{(1)}$	0.384(2)	0.320(1)	0.371(2)	0.308(2)	0.424(2)	0.350(1)
$\kappa_2^{(1)} \cdot \alpha_{2-2}^{(1)}$	−0.176(2)	−0.137(1)	−0.053(2)	−0.064(1)	−0.042(1)	−0.056(1)
Boosted						
$\kappa_1^{(1)} \cdot \alpha_{10}^{(1)}$	−0.228(4)	−0.281(7)	−0.201(4)	−0.297(7)	−0.017(1)	−0.037(1)
$\kappa_2^{(1)} \cdot \alpha_{20}^{(1)}$	0.165(8)	0.23(1)	0.123(8)	0.09(1)	0.648(2)	0.618(3)
$\kappa_2^{(1)} \cdot \alpha_{2-2}^{(1)}$	−0.005(7)	−0.07(1)	1.522(8)	1.19(1)	0.140(1)	0.154(2)

can be expected that α_{20} is more likely to receive larger QCD corrections than α_{10} . The QCD effects are much larger in the boosted setup rather than in the inclusive one, both in the SM and in the presence of SMEFT effects. For $m = 0$ coefficients, the inclusion of linear SMEFT terms has a minor impact compared to the quadratic EFT terms. This is somewhat expected as the quadratic term gives the leading change in the polarization balance [36], owing to the opening of helicity configurations that are suppressed in the SM [50,51]. The situation is dramatically different for the only nonvanishing azimuthal coefficient $\alpha_{2-2}^{(1)}$. In the inclusive setup, its value is strongly diminished by the SMEFT linear terms, while much milder effects come from the quadratic terms. In the boosted topology, this coefficient is extremely suppressed in the SM, while it

becomes positive owing to SMEFT contributions. While the SMEFT8 prediction is more or less of the same order of magnitude as in the inclusive setup, the SMEFT6 prediction is anomalously large, signaling an ill-defined extraction of azimuthal coefficients when the SMEFT series is truncated at dimension six. This goes in the same direction as the analytical conclusions of Sec. II. Indeed, the LO SMEFT6 prediction leads to a trace of the squared ρ_W matrix larger than 1 for $C_W/\Lambda^2 = 1 \text{ TeV}^{-2}$, driven by the large value of $\alpha_{2-2}^{(1)}$. At NLO QCD the violation of the trace bound is achieved for larger values of the WC, owing to QCD corrections that sizeably diminish the purity of the system. This can be observed in the top panels of Fig. 1. While in the inclusive setup the EFT effects are milder and not sufficient to spot the ill definition of the SMEFT6

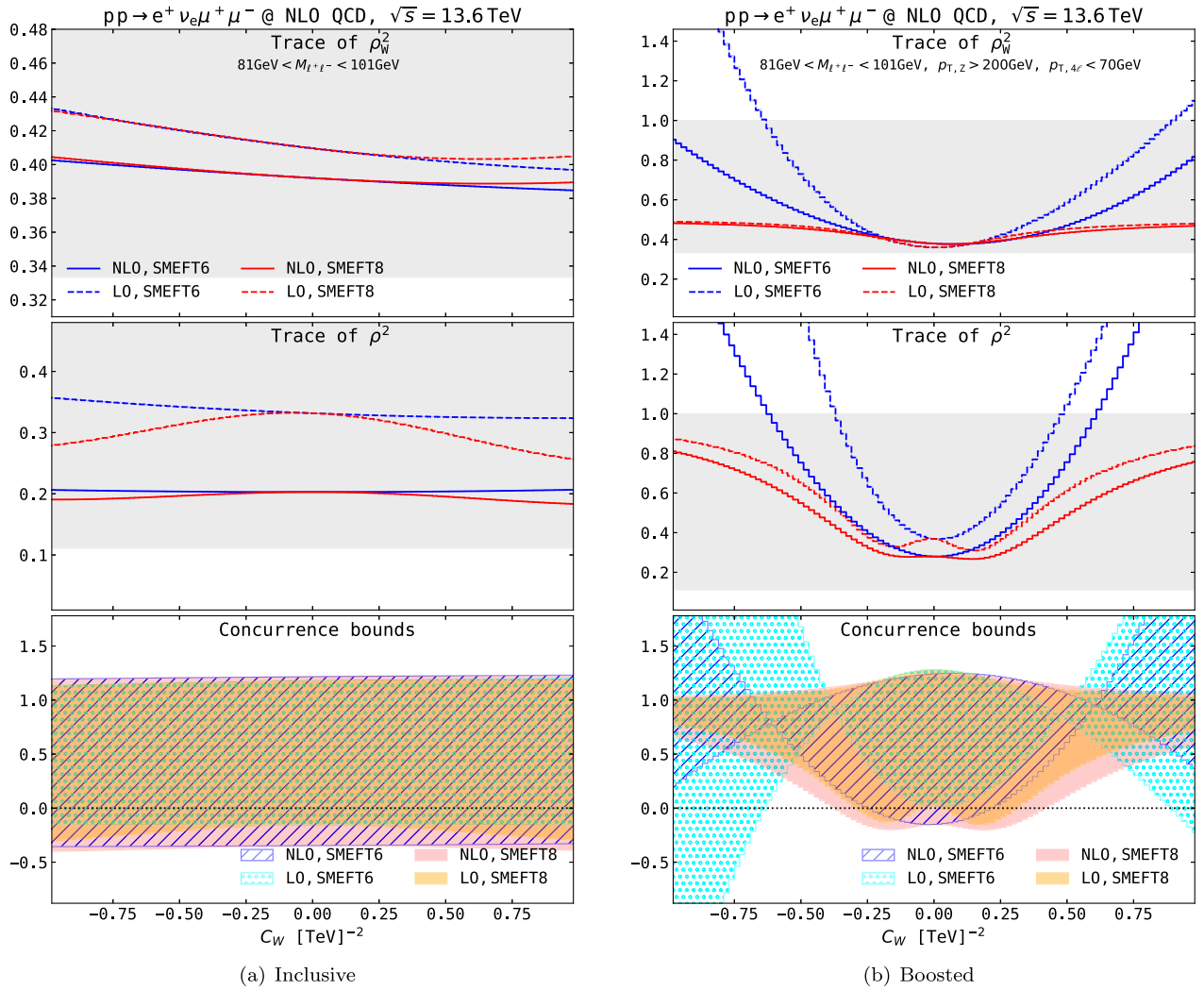


FIG. 1. Dependence of purity and concurrence markers on the Wilson coefficient C_W for inclusive WZ production at the LHC: the trace of the squared single-boson (W) spin-density matrix (upper panels), the trace of the squared ρ matrix associated to the overall WZ system (middle panels), and the concurrence lower and upper bounds for the WZ system (lower panels). Results have been computed at both LO and NLO QCD accuracy, including SMEFT effects up to linear (SMEFT6) and quadratic (SMEFT8) level. The inclusive (boosted) setup described in Eq. (26) [Eq. (27)] is considered in the left (right) plot. In the upper and middle panels, the gray bands correspond to allowed values for generic Hermitian, semipositive definite, unit-trace spin-density matrices.

predictions, in the boosted setup the sizeable EFT contributions expose in a striking manner the issues of Eq. (18). In kinematic topologies where the spin state has maximal purity, this effect would be even larger. From the top panel of Fig. 1(b) it is crystal clear that including quadratic terms is of paramount importance for a proper definition of the W-boson spin-density matrix, regularizing the ill-defined behavior of a ρ_W matrix as in Eq. (18). It is also important to notice that for large values of the WC and for boosted topologies the inclusion of NLO QCD corrections has a lower impact on SMEFT8 predictions than on SMEFT6 ones. As a last comment on single-boson coefficients, one could naively argue that the SMEFT6 predictions in the boosted topology allow one to constrain better the C_W value rather than using SMEFT8 ones. While tempting from the EFT-fit viewpoint, we refrain from drawing conclusions in this direction, not only because of the ill-defined spin-density matrix, but also because in this kinematic regime the SMEFT predictions start violating perturbative unitarity, and therefore one can expect that observables like the transverse momentum of the boson are expected to constrain better the WCs compared to angular coefficients.

Before discussing two-boson correlations, we point out that the results for the Z-boson spin-density matrix are very similar to those for the W boson in Table I and upper panels of Fig. 1, therefore we have not shown them explicitly in the text.

Looking at the red curves in the middle panels in Fig. 1, we observe that for a large WC value in the boosted topology the purity of the WZ-system spin state (≈ 0.8) is higher than the one of the W boson (≈ 0.5). Since the helicity rules select $W_{\pm}Z_{\pm}$ configurations (with the same share of left and right polarizations) and the size of single-boson azimuthal coefficients is moderate, the enhancement of the diboson purity only comes from diboson correlations and in particular from large values of the polar coefficient γ_{1010} , which depends on the relative share of left and right helicity of the two bosons [27], and of the azimuthal correlations $\gamma_{2\pm 22\mp 2}$. On the contrary, in the inclusive setup the purity of the WZ system (≈ 0.2) is lower than the one of the W (≈ 0.4), owing to very small azimuthal correlations. Overall, the inclusion of QCD corrections diminishes the purity of both the W and the WZ systems, as expected especially from the large real-radiation corrections that drive the NLO QCD correction even in the presence of constraints on the hadronic activity ($p_{T,WZ} < 70$ GeV). As discussed for ρ_W , the SMEFT6 description violates the constraint $\text{Tr}[\rho^2] \leq 1$ as well in the boosted regime. This violation appears for even smaller values of the C_W/Λ^2 than in the case of ρ_W , for both LO and NLO QCD predictions.

We have discussed at the end of Sec. II that the traces of the squared single- and diboson spin-density matrices can be combined into a lower and an upper bounds for the concurrence marker for spin entanglement. The

corresponding numerical results are presented in the lower panels of Fig. 1. Because of the violation of the $\text{Tr}[\rho^2] \leq 1$ and $\text{Tr}[\rho_V^2] \leq 1$ constraints in the SMEFT6 picture, the upper bound on the concurrence crosses the lower bound for absolute values of the Q_W -operator WC below 0.5. Even more strikingly than for the diboson spin purity, the SMEFT6 description of the spin density matrix fails in the boosted regime and for values of the WC that are still allowed by EFT fits. The inclusion of the quadratic terms in the SMEFT8 predictions regularizes the results. In the inclusive setup the SMEFT6 and SMEFT8 bounds look substantially the same and do not allow for any conclusion about spin entanglement, as the lower bound is negative. On the contrary, the SMEFT8 predictions in the boosted setup show a lower concurrence bound that exceeds zero for $|C_W/\Lambda^2| \approx 0.4(0.3)$ at NLO QCD (LO), highlighting a sizeable entanglement of the WZ spin state. As for the purity markers, the NLO QCD corrections diminish the lower bound of concurrence, pointing toward quantum-decoherence effects from QCD radiation.² We note that at LO in the boosted setup, while for values of the WC around 0.2 the lower bound is negative, in the strict SM limit it becomes slightly positive, in principle allowing for a claim of entanglement. However, introducing NLO QCD corrections makes the lower limit negative also for the SM. This means that, in spite of the possibility to optimize phase-space selections to enhance entanglement [23,24], the inclusion of higher-order QCD corrections could sizeably change the LO pictures, even in the presence of realistic jet vetoes.

The shape of the purity and concurrence bounds in the SMEFT8 predictions comes from the sums of squared α and γ coefficients in Eqs. (7)–(11). Each of the coefficients is the ratio of two quadratic polynomials in $1/\Lambda^2$ [see Eq. (15)]. For C_W different but close to zero, the numerator of several azimuthal coefficients takes a nonvanishing linear term, which is instead suppressed (due to helicity rules) in the denominator, leading to a concavity flip. On the other hand, for large values of C_W the quadratic terms become dominant in both the numerator and denominator, making the traces asymptotically approaching 1.

Although beyond the purposes of this work, we believe that the inclusion of spin correlations and entanglement markers as inputs to global SMEFT fits [83–85] is worth a try, even if it is not clear that their inclusion would improve the WC bounds. Recent single-operator sensitivity studies [62,64] suggest that polarization-sensitive observables can help constrain further SMEFT effects, compared to standard observables such as the WZ transverse mass. However, the practical utility of QI-inspired observables in SMEFT fits would only become clear after extending our study to multiple SMEFT operators. This should be done

²Quantum decoherence from radiative corrections at colliders has been recently investigated, e.g., in Refs. [80–82].

with care to avoid similar conceptual issues as those we discussed for SMEFT6 predictions in this work. For example, we checked numerically that the same issues with ρ positivity in the SMEFT6 predictions are present when activating two operators simultaneously (tested for Q_W and its CP -odd counterpart, $Q_{\bar{W}}$). This clearly shows that including multiple independent dimension-six operators in the definition of the spin-density matrix requires one to include not only quadratic terms for all operators, but also the interferences between them.

Before drawing our conclusions, we remark that the inclusion of dimension-8 operators is a natural extension of our work. It was recently shown [64] that genuine dimension-eight contributions to WZ from operators with maximal energy growth are subdominant compared to squared dimension-six effects from operators like Q_W . This holds for a subset of angular coefficients and for an even more boosted regime than the one considered in our study ($p_{T,V} > 400$ GeV). However, we refrain from excluding potentially sizeable effects from the interference of dimension-eight operators with SM amplitudes, which may have an impact on azimuthal coefficients and therefore could lead to a slightly modified picture for entanglement and purity markers.

On a similar note, we recall that NLO EW corrections in the SM [38,39] are about -5% in the boosted setup. While the effects of full NLO EW modeling in the SM heavily affects the Z-boson spin-analyzing power [32], we expect that the leading EW Sudakov logarithms do not sizeably impact the angular coefficients in the boosted setup, because of their factorized form as

$$d\sigma_{LL}^{\text{Sud}} \propto d\sigma_{LO} \alpha_{ew} \log^2 \frac{Q^2}{M_W^2} (+ \text{sub-leading}),$$

where Q is the invariant mass of the four-lepton system. Since angular coefficients are extracted integrating over Q as in Eq. (2), the impact of EW Sudakov logs should be minor. However, only a calculation of NLO EW corrections in SMEFT could unveil possible modifications to our NLO QCD picture for purity and entanglement markers. This is currently beyond reach.

V. CONCLUSION

We have presented a comprehensive analysis of spin correlations in the inclusive LHC production of WZ pairs in the fully leptonic channel. We have included NLO corrections in the strong coupling and new-physics contributions introduced by a SMEFT dimension-six operator that leads to anomalous electroweak TGCs. We have carried out the full quantum-state tomography of the considered diboson system, and connected the resulting entries of the spin-density matrix to standard markers of spin purity and entanglement.

We have shown that NLO QCD effects are very large on many polarization and spin-correlation coefficients, both in

the SM and in the SMEFT. Furthermore, the inclusion of higher-order QCD corrections was proved to diminish the level of purity and of entanglement of the two-boson spin state, signaling decoherence effects.

The insertion of CP -even SMEFT vertices modifying the TGCs has a moderate impact on angular coefficients in inclusive setups, while larger effects and nontrivial features appear for bosons with large transverse momenta.

Strikingly, in boosted topologies and for certain ranges of the considered SMEFT WC, the trace of the squared spin-density matrix exceeds one if the angular coefficients are extracted truncating the SMEFT expansion at dimension six, i.e., no quadratic EFT terms. Consequently, the naïve interpretation of angular coefficients in terms of the spin-density-matrix entries becomes cumbersome. A meaningful spin interpretation for arbitrary values of WCs and arbitrary kinematic regimes can be obtained only if quadratic EFT terms are properly included in the definition of the spin-density matrix and in the calculation of angular coefficients.

We also conclude that in realistic kinematic regimes, even boosted ones employed for recent LHC polarization measurements, as well as for realistic values of the SMEFT WC allowed by global fits, the concurrence bounds do not allow one to show that the WZ system is spin entangled. This is especially true after including NLO QCD corrections, while the LO results look more promising in this sense. Once again, we stress that the LO picture of spin correlations and entanglement markers is limited as far as the LHC production of boson pairs is concerned, making it vital to include at least NLO (and possibly next-to-next-to-leading order, which is now within reach) QCD corrections.

This study was focused on a single SMEFT operator modifying the TGCs for reasons of simplicity. However, we believe our conclusions would apply also to other operators, possibly affecting both the production and the decay of EW bosons.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [86].

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