New phenomenological and theoretical perspective on anomalous ZZ and $Z\gamma$ processes

Brando Bellazzini[1,*] and Francesco Riva[2,†]
[1]Institut de Physique Théorique, Université Paris Saclay, CEA, CNRS, F-91191 Gif-sur-Yvette, France
[2]Départment de Physique Théorique, Université de Genève, 24 quai Ernest-Ansermet, 1211 Genève 4, Switzerland

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Neutral diboson processes are precise probes of the Standard Model (SM) of particle physics, which entails high sensitivity to new physics effects. We identify in terms of dimension-8 effective operators the leading departures from the SM that survive in neutral diboson processes at high energy and that interfere with the unsuppressed SM helicity contributions. We describe symmetries and selection rules that single out those operators, both for weakly and strongly coupled physics beyond the SM. Finally, we show that unitarity and causality enforce, via dispersion relations, positivity constraints on the coefficients of these effective operators, reducing the parameter space which is theoretically allowed.

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I. MOTIVATION

Standard Model (SM) precision tests constitute an important part of the LHC physics program, in which SM predictions are confronted against precise data from proton-proton collisions. Beside testing our knowledge of SM processes, this program can be thought of as a searching tool for physics beyond the Standard Model (BSM). Indeed, new heavy particles beyond the energy reach of the collider might induce, via their virtual exchange, departures from the SM expectations. The well-established context in which SM precision tests are studied is that of effective field theories (EFTs), in which departures from the SM are organized as an expansion in inverse powers of a scale Λ, associated with the physical mass of putative new heavy resonances,

$$\mathcal{L}^{\rm eff} = \mathcal{L}^{\rm SM} + \sum_i c_i^{(6)} \frac{\mathcal{O}_i^{(6)}}{\Lambda^2} + \sum_i c_i^{(8)} \frac{\mathcal{O}_i^{(8)}}{\Lambda^4} + \cdots, \quad (1)$$

with $\mathcal{O}_i^{(d)}$ operators of increasing dimensionality d and dimensionless coefficients $c_i^{(d)}$.

In *most* cases, the leading departures from the scattering amplitudes predicted by the SM stem from $d=6$ operators, and the series can be truncated there. Natural exceptions come in two kinds. First, there can be BSM scenarios that imply selection rules (e.g., because of symmetries in the underlying theory), such that all $d=6$ operators vanish; in this case, the leading departures from the SM might arise at $d=8$ [1–3] or even higher [4]. Second, given that the number of $d=6$ operators is finite, there can be processes which do not receive any contributions at order $d=6$, but are affected only by operators of higher dimensionality. Neutral diboson processes $pp\to ZV$ $(V=Z,\gamma)$ [5,6] are one of the most interesting examples of the latter; they are in fact traditionally interpreted as measurements of neutral triple gauge couplings (nTGCs) [7], which correspond to $d=8$ operators in the EFT language [8].[1] There are other observables that are first modified at $d=8$ but are more difficult to identify [9,10] (some of these might in fact be the unique probes of a Higgs as a pseudo–Nambu-Goldstone boson [11]). Despite the high dimensionality of the operator, the coefficient $c_i^{(8)}$ can be sizeable in theories with a relatively strong underlying coupling, so much so as to partly alleviate the E/Λ suppression that accompanies the contribution of these operators to physical amplitudes.

*brando.bellazzini@ipht.fr
†francesco.riva@unige.ch

[1]There are also $d=6$ operators modifying these processes [for instance, operators of the form $(\bar\psi\gamma^\mu\psi)(iH^\dagger \overleftrightarrow{D}_\mu H)$], but their contribution is suppressed at high energy E by powers of m_Z/E; in addition, these operators are well constrained by resonant single-Z-boson processes (e.g., at the large electron positron collider), and their impact in $ZZ, Z\gamma$ is negligible.

For this reason, we can expect measurable effects even from these higher-dimensional operators.

In this work, we discuss the dimension-8 operators that affect ZZ and $Z\gamma$ processes and propose an innovative way of studying those, which is appealing both from an experimental and from a theoretical point of view. On the one hand, departures from the SM that are unsuppressed at high energy, and possibly interfere with the SM in simple analyses, have the highest chance of being detected; these are the effects that are most interesting from an experimental perspective. On the other hand, operators that appear in well-motivated and well-structured BSM scenarios are interesting from a more theoretical perspective as well, as they represent entire classes of theories. The most interesting operators are those that meet both these properties.

In what follows (Sec. III), we propose a set of operators of this kind. They modify ZZ and $Z\gamma$ amplitudes at $O(E^4/\Lambda^4)$ (that is the unsuppressed behavior expected from dimensional analysis), some of them contribute to amplitudes with diboson $+-/-+$ helicities (which happen to be the dominant configurations in the SM; see Sec. II), and, finally, most of these operators can be generated at tree level in models with spin-2 resonances or arise in scenarios with nonlinear supersymmetry, where the SM fermions are pseudo-Goldstini [2], as we discuss in Sec. IV.

An interesting and curious aspect of certain effective field theories is the notion of *positivity*, which follows, via dispersion relations due to analyticity (causality), Lorentz invariance, and locality, from the simple requirement that the underlying microscopic theory be unitary; see e.g., Refs. [12,13] and references therein.[2] These positivity bounds do not hold generically, without further assumptions (see Refs. [18–20]) for dimension-6 operators; for this reason, they have not received much attention in contemporary EFT LHC phenomenological studies. They do imply, however, strict positivity of certain coefficients of our dimension-8 operators, as we discuss in Sec. V. This model-independent reasoning will provide an important way to focus experimental searches to a smaller region of parameter space.

II. STANDARD MODEL ANATOMY

A thorough understanding of the SM amplitude is necessary to assess what is (and can be) measured at colliders. We discuss here the partonic $2\to 2$ amplitude, which is the target of the simplest inclusive analyses. The SM tree-level contribution to the $\bar\psi\psi\to ZZ, Z\gamma$ processes is characterised by a $t(u)$-channel singularity structure that projects on states of arbitrary angular momentum J and is dominated by the transverse-transverse (TT) $+-/-+$ helicity amplitudes. Final states with equal helicity $++/--$ are suppressed by m_Z^2/E^2 ($E=\sqrt{s}$) at high energies [21].[3]

The longitudinal + transverse (LT) configuration is instead always suppressed in the high-E limit by m_Z/E. This can be easily understood by noticing that, in the limit of vanishing Yukawas (which makes sense for the type of processes we are considering), a Z_2 symmetry $H\to -H$ characterizes the SM Lagrangian, implying that amplitudes with an odd number of scalars (that include the Higgs or the longitudinal components of vectors in the high-energy limit) must be suppressed by a vacuum expectation value v and, by dimensional analysis, lead to the above factor m_Z/E.

Finally, for ZZ the longitudinal-longitudinal (LL) helicity is very small in the tree-level SM, suppressed by m_Z^2/E^2. This follows from the $t(u)$-channel SM structure that characterizes tree-level ZZ production in the SM, where the direct coupling of scalars (equivalent to the longitudinal Z polarizations at high energy) to light quarks is suppressed by their small Yukawas.[4]

The amplitudes that do not vanish at high energy are

$$\mathcal{A}^{\rm SM}_{\bar\psi_+\psi_-\to Z^+Z^-} = 2g^2_{Z\psi_-}\tan\frac{\phi}{2}+\cdots, \tag{2}$$

$$\mathcal{A}^{\rm SM}_{\bar\psi_+\psi_-\to Z^+\gamma^-} = 2g_{Z\psi_-}g_{\gamma\psi}\tan\frac{\phi}{2}+\cdots, \tag{3}$$

where dots denote terms suppressed by powers of m_Z/E, ϕ is the angle between the momentum of the incoming $h=+1/2$ helicity fermion (or antifermion) and outgoing $h=+1$ helicity vector (that can be a Z or γ), and the subscript $\pm$ denotes the sign of the helicity. Here, $g_{Z\psi}=g(T^3_\psi-\sin^2\theta_W Q_\psi)/\cos^2\theta_W$, and $g_{\gamma\psi}=eQ_\psi$. Amplitudes with opposite vector helicity are simply related to those of Eq. (2) by

$$\mathcal{A}^{\rm SM}_{\bar\psi\psi\to Z^+V^-}(\phi)=\mathcal{A}^{\rm SM}_{\bar\psi\psi\to Z^-V^+}(\pi+\phi) \tag{4}$$

(with $V=Z,\gamma$) and similarly for $\psi_+\leftrightarrow\psi_-$ interchange, keeping in mind that the Z couplings to fermions are chiral $g_{Z\psi_+}\leftrightarrow g_{Z\psi_-}$.

A final important piece of information, is the fact that the largest SM amplitude is associated with left-handed initial state quarks, due to the known suppression of the right-handed quarks coupling to the Z boson: $\{g_{Zd_L},g_{Zu_L},g_{Zu_R},g_{Zd_R}\}\sim\{-0.32, 0.27, -0.10, 0.05\}$.

[2] For earlier applications in the QCD chiral Lagrangian, see e.g., Refs. [14–17].

[3] In the SM, dibosons are dominantly produced by quark-antiquark collisions with subleading quark-gluon and one-loop gluon-induced $gg\to VV$ components [22]; we do not consider these effects here.

[4] The situation is different at next-to-leading order, where a gluon-initiated, top-loop mediated diagram contributes sizeably to the LL final state.

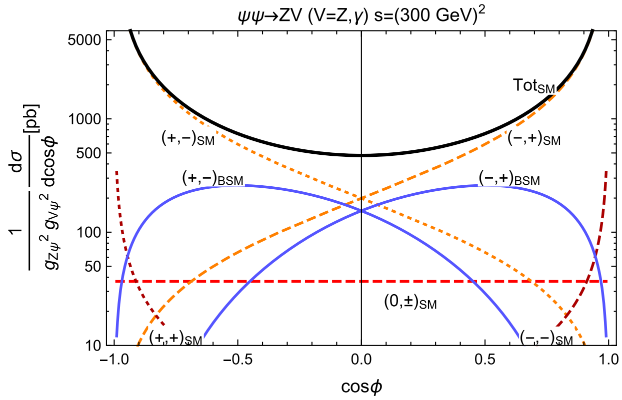

FIG. 1. Differential $\bar{\psi}\psi \to ZZ, Z\gamma$ cross section for different helicities, in units of the $\gamma\psi, Z\psi$ couplings, $g_{\gamma\psi}, g_{Z\psi}$ ($n_V = 1$ for $Z\gamma$ and $n_V = 2$ for ZZ). Dashed colored lines correspond to the SM-only contribution (solid black, the sum over helicities), while solid blue lines correspond to the BSM-only TT polarizations, with an arbitrarily chosen normalization, to be shown in the same plot as the SM.

The content of this section is summarized in Fig. 1. This can be trivially extended to $\bar{\psi}\psi \to \gamma\gamma$ processes, which we do not discuss in detail, as it is not traditionally discussed within the nTGCs framework.

III. EFFECTS BEYOND THE STANDARD MODEL

The nTGCs parametrization of Ref. [7], and its EFT counterpart [8],[5] is based on the physical hypothesis that the putative underlying new dynamics only modifies interactions among three gauge bosons. Consequently, new physics can enter $\bar{\psi}\psi \to ZV$ only through an s-channel diagram, or, more concretely, only the $J = 1$ amplitude can be modified. Because of angular momentum selection rules, only a longitudinal + transverse (LT) diboson final state is then possible for the neutral SM gauge bosons [23]. This amplitude is suppressed by m_Z/E also in BSM, and its interference with the SM is even more suppressed, since the SM LT piece is small.

Here, we take a step forward and ask ourselves how $\pm\mp$ diboson helicities can be sourced. The first configuration that allows for this is $J = 2$ and is sourced, for initial states involving fermions, by an operator containing $(i\bar{\psi}\gamma^{\{\mu}D^{\nu\}}\psi + \text{H.c.})$. From this, the leading CP even operators that we can write and lead to a $ZZ, Z\gamma$ final state are

$$\mathcal{O}^{(8)}_{\psi B} = -\frac{1}{4}(i\bar{\psi}\gamma^{\{\rho}D^{\nu\}}\psi + \text{H.c.})B_{\mu\nu}B^{\mu}_{\rho} \tag{5}$$

$$\mathcal{O}^{(8)}_{\psi W} = -\frac{1}{4}(i\bar{\psi}\gamma^{\{\rho}D^{\nu\}}\psi + \text{H.c.})W^{a}_{\mu\nu}W^{a\mu}_{\rho} \tag{6}$$

$$\mathcal{O}^{(8)}_{\psi H} = \frac{1}{2}(i\bar{\psi}\gamma^{\{\mu}D^{\nu\}}\psi + \text{H.c.})D_{\mu}H^{\dagger}D_{\nu}H \tag{7}$$

in the neutral channel and

$$\hat{\mathcal{O}}^{(8)}_{QH} = \frac{1}{2}(i\bar{Q}\sigma^{a}\gamma^{\{\mu}D^{\nu\}}Q + \text{H.c.})D_{\mu}H^{\dagger}\sigma^{a}D_{\nu}H \tag{8}$$

$$\hat{\mathcal{O}}^{(8)}_{QBW} = -\frac{1}{4}(i\bar{Q}\sigma^{a}\gamma^{\{\rho}D^{\nu\}}Q + \text{H.c.})B_{\mu\nu}W^{a\mu}_{\rho} \tag{9}$$

in the isospin-charged channel; here, $\psi = Q_L, u_R, d_R$ denotes the initial state fermion relevant for the LHC. These operators are dimension 8; that is, they are of the same order as effects that are tested at present as anomalous nTGCs [8].

The interactions in Eqs. (5), (6), (9) modify the TT amplitudes, contributing with different combinations to $Z_T Z_T$ as well as $Z_T\gamma$ and $\gamma\gamma$ final states. For example, from Eqs. (5), (6), at high energy, we find for $\psi = u_R$ or d_R

$$\mathcal{A}^{\text{BSM}}_{\bar{\psi}_+\psi_- \to Z^+V^-} = \frac{c^{\text{tot}}_V}{4}\frac{\hat{s}^2}{\Lambda^4}\sin\phi(1+\cos\phi), \tag{10}$$

where $c^{\text{tot}}_Z = \sin^2\theta_W c^{(8)}_{\psi B} + \cos^2\theta_W c^{(8)}_{\psi W}$ and $c^{\text{tot}}_\gamma = \sin(2\theta_W)(c^{(8)}_{\psi B} - c^{(8)}_{\psi W})$ are the effective combinations that enter in the two processes (θ_W the weak mixing angle). Here, $\hat{s}$ is the center-of-mass energy, and Eq. (10) exhibits the unsuppressed energy growth that one expects from dimension-8 effects. The associated differential cross sections are shown in Fig. 1. Analogous expressions hold for initial states with $\psi = Q$, and $\gamma\gamma$ final states.

Similarly, Eqs. (7) and (8) contribute at high energy $E \gg m_Z$ to the production of two longitudinally polarized Z bosons, as can be understood by the equivalence theorem and dimensional analysis. From Eq. (7), we find at high energy

$$\mathcal{A}^{\text{BSM}}_{\bar{\psi}_+\psi_- \to Z_L Z_L} = \frac{c^{(8)}_{\Psi H}}{8}\frac{\hat{s}^2}{\Lambda^4}\sin(2\phi). \tag{11}$$

An important aspect of these precision SM tests is the interference between the SM process and the BSM effect: a sizeable interference enhances the sensitivity to these BSM effects. As discussed in the previous section, the majority of SM processes that we observe at the LHC have $\pm\mp$ helicity and left-handed initial state quarks. This amplitude is modified by the operators of Eqs. (5), (6), and (9) with $\psi = Q_L$. Furthermore, as can be observed from Fig. 1, the SM and BSM same-helicity polar-angle distributions have

[5]Notice that anomalous couplings always induce an energy growth in some scattering process, which ultimately implies a cutoff at energy $E \sim \Lambda^{\text{cutoff}}$, above which the associated theory ceases making sense; for this reason, anomalous coupling parametrizations can always be reformulated as EFTs with $\Lambda \leq \Lambda^{\text{cutoff}}$.

an important overlap, implying that a sizeable SM-BSM interference can be expected in fixed energy bins. Simple analysis with standard selection criteria, including a small cut in the very-forward region $|\cos\phi| \approx 1$, shall be already particularly sensitive to the deformations of Eqs. (5), (6), and (9).

Interference of the LL channel is instead suppressed, given the smallness of the LL SM amplitude (see the previous section). So, this channel maintains its interest mainly from its very-high-energy behavior and the BSM connection with modified Higgs dynamics.

On the other hand, it is interesting to notice that analyses based on the LT configurations (as implicitly assumed by nTGCs analyses) imply that this majority of SM events plays effectively the role of background, rather than signal. Instead, the operators of Eqs. (5), (6), and (9) source the right helicity amplitudes so that the situation is reversed and all available experimental information is systematically used and tested.

The final states we are interested in are not unique of the $J = 2$ angular momentum configuration and are also found for $J \geq 3$. However, states with larger angular momentum are necessarily sourced by local operators with additional powers of momenta/derivatives. That is, $J \geq 3$ is associated with operators of dimension > 8 that are, in this situation, negligible.

IV. EXAMPLE BSM MODELS

In the previous sections, we have motivated, from an experimental (bottom-up) point of view, a new class of effects that appears at the same order in the EFT expansion as effects in the new class are easier to detect (for a comparable new physics scale) given their unsuppressed nature and their interference with the SM. Now, we take a more theoretical (top-bottom) perspective and argue that these operators are also well motivated from a BSM point of view; we discuss two BSM scenarios in which these effects could arise with sizeable coefficients.

Here, it is perhaps worth pausing a moment to understand whether dimension-8 operators can be relevant at all, given that they appear to be subleading in the energy expansion. A given BSM that generates sizeable $d = 6$ effects as well as $d = 8$ ones will probably be better searched through those $d = 6$ effects. So, we are interested in whether it is possible that $c^{(6)} \ll c^{(8)}$, that is, situations where the coefficients of $d = 6$ operators are suppressed while $d = 8$ are not. This is certainly possible in some finely tuned region of parameter space, but this goes against the perspective that EFTs capture broad BSM scenarios, rather than specific points in parameter space. Symmetries and other selection rules can instead induce natural situations for this hierarchy, as we now discuss with two examples.

A first situation that leads to dimension-8 effects that are larger than dimension-6 ones is the tree-level exchange of weakly coupled massive spin-2 resonances, such as graviton Kaluza-Klein excitations in models with extra dimensions [24–26]. The massive graviton interacts (like its massless version) with the stress-energy tensor $T_{\mu\nu}$; its Lagrangian is

$$\mathcal{L}_g = -\frac{m_g^2}{2} h^{\mu\nu} P_{\mu\nu\rho\sigma} h^{\rho\sigma} - \frac{1}{\bar{M}_p} h^{\mu\nu} T_{\mu\nu} \tag{12}$$

with $P_{\mu\nu\rho\sigma} = (\eta_{\mu\rho}\eta_{\nu\sigma} + \eta_{\mu\sigma}\eta_{\nu\rho})/2 - \eta_{\mu\nu}\eta_{\rho\sigma}/3 + \cdots$, which is equivalent to the propagator expanded at leading order in momentum over the spin-2 mass p/m_g, and $\bar{M}_p$ the reduced Planck mass in the extra dimension. Integrating out h, one finds at leading order in $1/m_g$

$$\mathcal{L}_g^{\text{eff}} = \frac{1}{2m_g^2 \bar{M}_p^2} \left[(T^{\mu\nu} T_{\mu\nu}) - \frac{1}{3} (T^{\mu}_{\mu})^2 \right] + \cdots. \tag{13}$$

The second piece only leads to effects with off-shell fermions and is not relevant for our discussion (in fact, via a field redefinition, it can be written as a dimension-10 effect), while the first one leads to

$$c_{\psi H} = c_{\psi B} = c_{\psi W} = \frac{m_g^2}{\bar{M}_p^2} \tag{14}$$

with $\Lambda = m_g$, in addition to a number of other $d = 8$ operators involving four fermions, vectors, or scalars [24].

We have illustrated a model that singles out $d = 8$ operators, and in particular the ones proposed in this paper. This holds only in the limit where m_g is much lighter than other BSM resonances and moreover is weakly coupled, meaning that the relevant coupling at the scale m_g is $\ll 4\pi$.[6] Indeed, for strong coupling, graviton loops generate $d = 6$ effects [27], and, moreover, one expects a richer spectrum of states at the cutoff.

Symmetries provide instead a more robust context to study the relevance of $d = 8$ effects, which can also hold in the strongly coupled limit. Consider for instance a real scalar field ϕ, Nambu-Goldstone boson of a $U(1)$ symmetry, spontaneously broken at the physical scale Λ. The $U(1)$ symmetry translates into a shift symmetry $\phi \to \phi + \alpha$ [α a global $U(1)$ phase[, which implies that the effective Lagrangian respecting this (nonlinearly realized) symmetry can only involve powers of $\partial\phi$:

[6]Notice that most LHC processes are at present tested with large statistical uncertainty at high energy, aiming at deviations from the SM which are larger than $O(1)$; in this situation, a consistent EFT interpretation is possible only if the underlying theory is coupled more strongly than the SM [3]. This suggest a sizeable window of parameters where these models are represented by a $d = 8$ EFT and are consistently testable.

$$\mathcal{L}_{\phi}^{\rm eff} = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi + \frac{c_\phi}{\Lambda^4}(\partial_\mu\phi\partial^\mu\phi)^2 + \cdots \qquad (15)$$

Clearly, in this theory, it is natural that the leading interactions are dimension 8, a statement that is independent of the coupling strength in the microscopic theory, since the symmetry always protects against the generation of $d=6$ terms. This remains true even if the $U(1)$ symmetry is only approximate, broken explicitly by small parameters, such as a mass term $m_\phi^2\phi^2/2$, with $m_\phi \ll \Lambda$. In this situation, dimension-6 operators [such as $(\partial_\mu|\phi|^2)^2$] will be generated, but we can expect their coefficients to be small and controlled by the small parameter $m_\phi^2/\Lambda^2 \ll 1$.

The very same reasoning has been applied to the SM fields in Refs. [1,2,28]. In particular Ref. [2] discusses the analog of the $\phi \to \phi + \alpha$ shift symmetry in the context of fermions: nonlinearly realized extended supersymmetry. If the SM fermions are (pseudo-)Goldstini of spontaneously broken (SB) supersymmetry, then their leading interactions arise indeed at $d=8$ and provide an example where dimension-8 effects can be well motivated from a BSM perspective.

The low-energy physics of pseudo-Goldstini can be captured by an effective theory for SB space-time symmetries so that the main couplings between pseudo-Goldstini χ and matter arise effectively as a distortion of the effective metric perceived by matter fields in the SB background,

$$g^{\mu\nu} = \eta^{\mu\nu} + \frac{1}{2F^2}(i\bar{\chi}\gamma^\mu\partial^\nu\chi + i\bar{\chi}\gamma^\nu\partial^\mu\chi + \text{H.c.}) + \cdots, \qquad (16)$$

where F is the supersymmetry breaking scale. In this context, if the SM fermions are pseudo-Goldstini $\psi = \chi$, then the kinetic term for Higgs $D_\mu H^\dagger D_\nu H g^{\mu\nu}$ leads to Eq. (7), while the $V = W, B$ kinetic terms $-1/4 V^A_{\mu\nu} V^A_{\rho\sigma} g^{\mu\rho} g^{\nu\sigma}$ lead to the interactions of Eqs. (5) and (6), with

$$c_{\psi H} = c_{\psi B} = c_{\psi W} = \frac{\Lambda^4}{F^2}, \qquad (17)$$

with Λ the mass of the other supersymmetric particles.

The complete set of (approximate) symmetries that can lead to $c^{(8)} \gg c^{(6)}$ has been discussed in Refs. [1,2]: while nonlinear supersymmetry can protect Eqs. (6) and (7), there is no known analog for other operators entering diboson pair production. This puts the operators (6) and (7) on privileged ground, also from a BSM perspective.

V. POSITIVITY CONSTRAINTS AND BEYOND

Dispersion relations, following from the fundamental principles of causality (analyticity), locality (Froissart bound [29,30]), and crossing symmetry, imply relations between low energies (IR), where LHC experiments are performed, and high energies (UV), where the SM and the effects, Eqs. (5)–(9), are completed into a microscopic unitary theory. This implies strict positivity constraints for the coefficient of the s^2 term in the Taylor expansion of elastic scattering amplitudes at $t, s \to 0$,

$$\partial^2\mathcal{A}/\partial s^2|_{s,t=0} > 0. \qquad (18)$$

This has important consequences for EFTs with unknown, but unitary, UV completions [12,13]. By dimensional analysis, it is clear that this is a unique feature of operators of dimension $d=8$ [or $d \geq 8$, via higher derivatives in Eq. (18)] that is unparalleled in the more familiar framework of $d=6$ operators (see, however, Refs. [18–20] for a similar relation involving additional inputs from the UV).

Applied to the operators of Eqs. (5)–(7), appearing in the Lagrangian as in Eq. (1), given that $\Lambda^4 > 0$, we find[7]

$$c_{\psi H}^{(8)} > 0, \qquad c_Z^{\rm tot} > 0, \qquad c_{\gamma\gamma}^{\rm tot} > 0 \qquad (19)$$

for $\psi = u_R, d_R$, where $c_{\gamma\gamma}^{\rm tot} = c_{\psi B}^{(8)}\cos^2\theta_W + c_{\psi W}^{(8)}\sin^2\theta_W$. That is, fundamental principles reduce the parameter space that can be explored to half its size and help focus experimental searches. These (strictly positive) constraints also imply that these operators are necessarily there, and it would not make sense to study, for instance, scenarios where BSM starts with leading $d > 8$ operators (or $J > 3$ angular momentum) [13]. Scattering instead $\psi = Q = (u_L, d_L)$ states, the operators in (8) and (9) contribute to the forward elastic amplitudes as well, the results varying with the isospin $\sigma^3_{ii} = \pm 1$ associated to u_L and d_L, respectively,

$$\mathcal{A}_{QZ_L \to QZ_L}(s, t=0) = (c_{QH}^{(8)} \mp \hat{c}_{QH}^{(8)})\frac{s^2}{\Lambda^4}, \qquad (20)$$

$$\mathcal{A}_{QV \to QV}(s, t=0) = c_{QV}^{(8)\pm}\frac{s^2}{\Lambda^4}. \qquad (21)$$

Here, $V = Z_T, \gamma$, and we defined $c_{QZ}^{(8)\pm} = \sin^2\theta_W c_{QB}^{(8)} + \cos^2\theta_W c_{QW}^{(8)} \mp \hat{c}_{QBW}^{(8)}\sin(2\theta_W)/2$, $c_{Q\gamma}^{(8)\pm} = c_{QB}^{(8)}\cos^2\theta_W + c_{QW}^{(8)}\sin^2\theta_W \pm \hat{c}_{QBW}^{(8)}\sin(2\theta_W)/2$. Therefore, positivity implies that

[7] The following positivity bounds are extracted at very small energy, where the Wilson coefficients differ from those at the scale Λ [e.g., Eqs. (14) and (17)]. Here, we assume that the rest of the theory (i.e., the SM and possible dimension-6 operators) is weakly coupled below Λ, so that (calculable) renormalisation group evolution effects can be neglected, as long as $c^{(8)}$ is not much larger than $(4\pi)^2$, because otherwise insertions of operators with $d < 8$ and at least one dimension-8 operator with such a large coefficient might generate large radiative corrections.

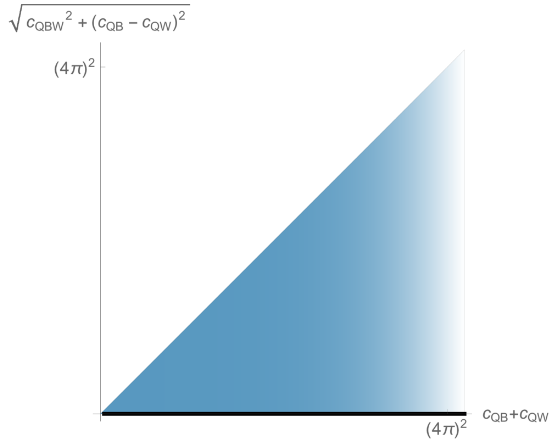

FIG. 2. Parameter space for the operators with $\psi = Q$ (we have neglected superscripts for clarity). The blue region is allowed by positivity constraints, Eqs. (23) and (24). The black segment corresponds to the explicit scenarios of Sec. IV.

$$c_{QH}^{(8)} \mp \hat{c}_{QH}^{(8)} > 0, \qquad c_{QV}^{(8)\pm} > 0 \tag{22}$$

and that $\hat{c}_{QBW}^{(8)}$ cannot appear without (equally large) contributions from $c_{QV}^{(8)\pm}$.

The statement of Eq. (22) can be made even stronger if the dynamics that generates these operators is insensitive to the low-energy electroweak and electroweak symmetry breaking physics, that is, if the new sector admits a consistent UV completion irrespectively of the Weinberg mixing angle θ_W which is determined by the external weak gauge couplings.[8] More explicitly, we are advocating for the gedanken limit $g \to 0$ and $g' \to 0$ while allowing the ratio $g'/g = \tan\theta_W$ to be free and independent of the strong sector. This means that the bounds (19) and (22) hold for any θ_W, which in turn implies that

$$c_{\psi H}^{(8)} > 0, \qquad c_{\psi W}^{(8)} > 0, \qquad c_{\psi B}^{(8)} > 0, \qquad c_{QH}^{(8)} > |\hat{c}_{QH}^{(8)}|, \tag{23}$$

with $\psi = Q$, u_R, and d_R, and additionally

$$4 c_{QB}^{(8)} c_{QW}^{(8)} > (\hat{c}_{QBW}^{(8)})^2. \tag{24}$$

The inequalities for $\psi = Q$, which derive from Eq. (21) using θ_W as a free parameter, define a one-branch cone in the $(c_{QB}^{(8)}, c_{QW}^{(8)}, \hat{c}_{QBW}^{(8)})$ space, which is obtained by revolving the triangle shown in Fig. 2 around the axis $c_{QB}^{(8)} + c_{QW}^{(8)}$. These generic results are of course compatible with the specific models discussed above, Eqs. (14) and (17), where $\hat{c}_{QBW}^{(8)}$ and $\hat{c}_{QH}^{(8)}$ are not generated, and all other operators are generated with positive coefficients.

In addition to the above-mentioned *lower* bounds on the coefficients $c^{(8)}$, there are various arguments that suggest *upper* bounds. These arguments include perturbativity, beyond positivity [31], or simple naive dimensional analysis [32,33] and indicate that for coefficients $c \gg (4\pi)^2$ the EFT description that we have proposed (where only $d = 8$ operators are retained) may break down.

VI. CONCLUSIONS AND OUTLOOK

In this work, we have provided a parametrization for neutral diboson processes with initial state fermions, which extends the traditional parametrization based on neutral triple gauge couplings. The scope has been to provide a framework for testing the precise SM predictions [34–39], against the most interesting and well-motivated alternative hypotheses. Our arguments have been exposed in terms of an effective field theory (where they are naively of the same order as nTGCs effects) but are based on the general helicity properties of diboson amplitudes. In particular, we point out $J = 2$ effects that are unsuppressed at high energy, both in the LL and TT final states, and highlight effects that do interfere with the dominant SM amplitude which have $\pm\mp$ helicity. In this sense, our proposal goes toward the goal of exploiting at best LHC data, providing a way to present information about the high-energy SM amplitudes in its completeness, and identifying the features that can be tested most precisely.

We then approached these operators from a more theoretical perspective and argued that fundamental principles based on unitarity imply generic and model-independent constraints on the operator coefficients. This is *per se* a very interesting result that has no general analog in the more familiar context of $d = 6$ operators. As a result, the parameter space of the naive EFT is drastically reduced, thus focusing experimental efforts to the relevant physical scenarios. We find for instance that the operators of Eqs. (8) and (9), when taken in isolation, cannot be completed into a unitary theory, despite being allowed from a naive EFT point of view.

Finally, we motivated the proposed effects from a BSM point of view. We have discussed a specific model where the virtual tree-level exchange a massive spin-2 resonance (KK graviton) generates the operators under scrutiny. In this model, $d = 6$ operators are loop-suppressed, while $d = 8$ operators are not. Therefore, in the weakly coupled regime, this example produces its larger effects in our $d = 8$ operators. Moreover, we have identified a symmetry (nonlinearly realized supersymmetry) that singles out the operators, Eqs. (6) and (7), and holds also in the strongly coupled regime. This puts the hierarchy $c^{(8)} \gg c^{(6)}$ of our scenario on firm ground in the realm of both weakly and strongly coupled models.

[8] Examples of this are the Kaluza-Klein graviton or Goldstini of Sec. IV, where the transverse gauge bosons are also part of the strongly interacting sector, as described in Ref. [1].

This opens several interesting prospects for future research. From an experimental point of view, it will be interesting to asses the reach of LHC experiments to this type of physics [in particular to the operators of Eqs. (6) and (7) with $\psi = Q_L$], with present and future data, and eventually implement this strategy into present ZZ and $Z\gamma$ studies. Contrary to charged diboson processes, $ZZ, Z\gamma$ give easy access to the center-of-mass energy $\hat{s}$, facilitating therefore a discussion of the EFT validity, along the lines of Ref. [3]. Moreover, resonance searches for spin-2 (Kaluza-Klein graviton) states are already performed at the LHC in the ZZ, $Z\gamma$, or $\gamma\gamma$ channels [40–43]; it would be nice to compare and combine these complementary modes of exploration (resonant and EFT) into a single unified picture. Finally, it would be interesting to discuss next-to-leading-order effects and in particular their contribution to the 00 amplitude to understand whether this can be exploited to search for specific BSM scenarios, e.g., Refs. [44,45], in the $gg \to ZZ$ amplitude.

From a model-building point of view, it would be interesting to put the pseudo-Goldstino models of Ref. [2] on firmer ground, stressing their relation to the hierarchy problem and providing additional motivation for these searches. Finally, the reasoning we develop here could be extended to charged diboson final states $W^\pm Z$, W^+W^-. Indeed, it has been shown on completely general grounds that $d = 6$ effects including transverse vectors have suppressed SM-BSM interference in the high-energy limit [21] for inclusive searches (see, however, Refs. [46,47]); in these conditions, the effects discussed here and those of the dimension-6 operator W^3 would appear at the same order, a feature that might deserve a dedicated discussion in explicit scenarios.

Our proposal offers an innovative opportunity for experiments transiting toward an EFT parametrization of non-SM effects.

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