Vices and virtues of Higgs effective field theories at large energy

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We study constraints on new physics from Higgs production at the LHC in the context of an effective field theory (EFT), focusing on Higgs searches in HV ($V = W, Z$) associated production which are particularly sensitive to the high-energy behavior of certain dimension-six operators. We show that analyses of these searches are generally dominated by a kinematic region where the generic EFT expansion breaks down, and we establish under which conditions they can nevertheless be meaningful. For example, constraints from these searches on the Wilson coefficients of operators whose effects grow with energy can be established in scenarios where a particular combination of fermions and the Higgs are composite and strongly coupled: then, bounds from Higgs physics at high energy are complementary to LEP1 and competitive with LEP2.

I. MOTIVATION

With the discovery of a Higgs boson [1,2], experiments have finally probed all sectors of the Standard Model (SM). The priority is now to measure the properties of the Higgs particle and to explore the mechanism of electroweak symmetry breaking. In this paper, we will focus on scenarios beyond the SM (BSM) in which the new physics modifying the Higgs interactions is heavy.[1] Its leading effects on the SM can then be parametrized through effective operators of dimension six, suppressed by the scale of new physics Λ [6,7]. Some of the many possible operators which affect Higgs properties can be measured in Higgs physics only, while others are related to electroweak (EW) observables. This is due to the fact that the Higgs scalar excitation $v + h$ is always associated with the EW symmetry breaking order parameter v [8–10]. Due to the limited precision of hadron machines, one would think that LHC Higgs measurements are unlikely to compete with LEP constraints on this second group of operators. Nevertheless, the extended energy reach of LHC allows it to access regions where the effects of some operators are enhanced by powers of E/Λ, leading to an increase in sensitivity.

Unlike in on-shell Higgs production by gluon fusion or in Higgs decays, which occur at $E \sim m_h$, in channels in which the Higgs is produced in association with electroweak gauge bosons, $pp \rightarrow hV$, $V = W, Z$, the invariant mass flowing into the hVV vertex is mainly limited by PDF suppression: these channels can have enhanced sensitivity to effects growing with energy. We quantify the extent to which the corresponding cross sections and kinematic distributions [11–13] can be used to constrain physics beyond the SM. We find that the naive bounds on the coefficients of dimension-six operators which can be extracted from these measurements are indeed very strong, even with the limited amount of data available at present. Nevertheless, the E/Λ enhancement comes at the cost that these measurements are dominated by kinematic regions where the effective field theory (EFT) expansion has broken down unless specific Wilson coefficients are very large. They are, therefore, meaningless in the context of generic EFTs. Our arguments hold also, e.g., for analyses of the high-energy tail of the $pp \rightarrow h^* \rightarrow ZZ$ cross section [14–18] and can be used also in this context to estimate the regime of validity of the EFT.

The very motivation for studying EFTs and their Wilson coefficients is that they allow for a simple parametrization of experimental constraints and an efficient comparison with large classes of UV theories. For this reason it is important to attribute a physical meaning to the Wilson coefficients in terms of masses, couplings or multiplicities of the BSM sector. It is, thus, crucial to understand which classes of theories (if any) can yield and enhancement of Wilson coefficients contributing to Higgs physics at high energy, rather than simply assuming that such theories exist. Only then can the bounds which we extract be thought to carry some information.

With this motivation in mind, we assume that the underlying new physics is under perturbative control even in the strong coupling limit where it can be thought of as the effective description of a composite sector. We then integrate out minimally coupled massive states to match this sector to the SM EFT description. We show that one

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[1]Scenarios with additional light non-SM particles are either tightly constrained already or can be searched for through their modifications of differential distributions at small momenta [3–5].

combination of operators, $\mathcal{O}_W - \mathcal{O}_B$, which contributes dominantly to $pp \to hV$ production (but not at tree level to other tightly constrained observables such as $h \to \gamma\gamma$, $Z\gamma$ or EWPTs [9]) can indeed be enhanced by a strong coupling in the underlying theory if the Higgs and a particular combination of fermions are composite and strongly coupled [19]. In this context, the bounds derived from hV associated production are surprisingly strong, complementary to those from LEP1 and competitive with those from LEP2 (see also [20,21]). In theories where our bounds are consistent (and, for instruction, also in theories where they are not), we compare our results with the bounds from LEP2 measurements of Triple Gauge Couplings (TGCs), which receive contributions from the same operators.

This article is organized as follows: In Sec. II we introduce the effective field theory description of Higgs physics. We choose a basis of operators which is particularly well-suited for our needs, as it not only allows a relatively straightforward interpretation in terms of observables, but can also be easily matched to relevant models of underlying new physics. We discuss the connection between $pp \to hV$ and observables in TGCs and examine the high-energy behavior of the operators in question. The validity of the effective field theory description is examined in Sec. III where we consider explicit models of new physics with heavy vector resonances. In Sec. IV we analyze the existing data for associated Higgs production with W or Z bosons in ATLAS. Informed by our discussion of EFT validity and breakdown, we study which bounds on the coefficients of the higher dimensional operators can be established under different assumptions about the underlying new physics. We conclude in Sec. V.

II. DIMENSION-SIX OPERATORS IN HIGGS PHYSICS

The lack of direct discovery of BSM physics suggests that, if such physics exists, it is much heavier than the EW scale and lies beyond the LHC reach. In this situation new physics (NP) can still leave an indirect imprint in low-energy observables. This can be efficiently and generically parametrized in the context of EFTs, corresponding to an expansion in the SM fields and derivatives over the NP scale Λ,

$$\mathcal{L}_{\mathrm{eff}} = \mathcal{L}_4 + \mathcal{L}_6 + \cdots, \tag{1}$$

where $\mathcal{L}_4$ defines the SM, while $\mathcal{L}_6$ can be written as a sum of local dimension-six operators,

$$\mathcal{L}_6 = \sum_{i_1} g_*^2 \frac{c_{i_1}}{\Lambda^2} \mathcal{O}_{i_1} + \sum_{i_2} \frac{c_{i_2}}{\Lambda^2} \mathcal{O}_{i_2}, \tag{2}$$

where we have differentiated between two classes of operators [10]. The operators $\mathcal{O}_{i_1}$ involve extra powers of SM fields and for this reason must also involve extra

powers of a coupling which we have denoted generically by g_*.[2] The operators $\mathcal{O}_{i_2}$, on the other hand, involve extra powers of derivatives and are, thus, suppressed by the scale Λ only. In this notation, the Wilson coefficients c_i are dimensionless, both in mass and coupling units. If the expansion is valid, $\mathcal{L}_6$ parametrizes the dominant contributions of baryon- and lepton-number preserving NP. Complete sets of dimension-six operators can be found in Refs. [6,7,10,22], expressed in different bases, equivalent up to field redefinitions.

However, only a few of these operators contribute to Higgs physics, and some of them to EW physics as well. Hence, a global fit including all operators and all experiments becomes necessary [8,23,24]. This cumbersome task can be partially avoided (and much physical insight gained) with an educated choice of basis [9,25]. Ideally, it should allow us to identify exactly which operators are tightly constrained by LEP1 experiments (and can, therefore, be neglected in LHC physics) and which ones could still provide measurable deviations from the SM.

The most appropriate basis for our purposes, which describes the relevant contributions to Higgs physics, shares part of the advantages described in Refs. [9,25] but can also be quickly matched to specific UV models, is shown in Table I. We restrict this discussion to *CP*-even operators at the leading order in the minimal-flavor-violation hypothesis [26], but both assumptions can be easily relaxed. The bounds we derive can be quickly translated into other bases using operator identities such as

$$\mathcal{O}_B = \mathcal{O}_{HB} + \mathcal{O}_{BB} + \mathcal{O}_{WB},$$
$$\mathcal{O}_W = \mathcal{O}_{HW} + \mathcal{O}_{WW} + \mathcal{O}_{WB}, \tag{3}$$

where

$$\mathcal{O}_{HW} = ig(D^\mu H)^\dagger \sigma^a (D^\nu H) W_{\mu\nu}^a,$$
$$\mathcal{O}_{WB} = \frac{gg'}{4} (H^\dagger \sigma^a H) W_{\mu\nu}^a B^{\mu\nu}, \tag{4}$$

and field redefinitions proportional to the equations of motion (EOM),

$$\mathcal{O}_W = g^2 \left[\frac{3}{2} \mathcal{O}_r - \frac{1}{4} \sum_{u,d,e} \mathcal{O}_y - \mathcal{O}_6 + \frac{1}{4} (\mathcal{O}'_{HL} + \mathcal{O}'_{HQ}) \right].$$

$$\mathcal{O}_B = g'^2 \left[-\frac{1}{2} \mathcal{O}_T + \frac{1}{2} \sum_F Y_F \mathcal{O}_{HF} \right], \tag{5}$$

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[2]This is most clearly seen by keeping powers of $\hbar$ explicit in the action $S = \int d^4x \mathcal{L}/\hbar$; then, since a simultaneous rescaling of $\hbar$ and all the couplings and fields in $\mathcal{L}$ cannot modify S, the couplings must scale like $\hbar^{-1/2}$ while the fields scale like $\hbar^{1/2}$: this fixes the coupling-power counting of dimension-six operators.

TABLE I.   Complete, nonredundant list of *CP*-even dimension-six operators that can potentially contribute to Higgs physics. On the left, operators that can only affect Higgs physics [8–10]; on the right, operators already constrained by EW tests. Flavor indices are summed over (e.g. $\bar{e}_R\gamma^\mu e_R$ stands for $\bar{e}_R\gamma^\mu e_R + \bar{\mu}_R\gamma^\mu \mu_R + \bar{\tau}_R\gamma^\mu \tau_R$). For each operator, we indicate whether it belongs to class 1 or class 2 in the classification of Eq. (2) [10]. Our normalization of the operators differs from previous literature.

Higgs physics only	
$\mathcal{O}_r = |H|^2|D^\mu H|^2$	1
$\mathcal{O}_{BB} = \frac{g'^2}{4}|H|^2 B_{\mu\nu}B^{\mu\nu}$	2
$\mathcal{O}_{WW} = \frac{g^2}{4}|H|^2 W_{\mu\nu}^a W^{a\mu\nu}$	2
$\mathcal{O}_{GG} = \frac{g_s^2}{4}|H|^2 G_{\mu\nu}^A G^{A\mu\nu}$	2
$\mathcal{O}_{y_u} = y_u|H|^2 \bar{Q}_L \tilde{H} u_R$	1
$\mathcal{O}_{y_d} = y_d|H|^2 \bar{Q}_L H d_R$	1
$\mathcal{O}_{y_e} = y_e|H|^2 \bar{L}_L H e_R$	1
$\mathcal{O}_6 = \lambda|H|^6$	1

EW and Higgs physics	
$\mathcal{O}_W = \frac{ig}{2}(H^\dagger \sigma^a \overset{\leftrightarrow}{D^\mu} H) D^\nu W_{\mu\nu}^a$	2
$\mathcal{O}_B = \frac{ig'}{2}(H^\dagger \overset{\leftrightarrow}{D^\mu} H)\partial^\nu B_{\mu\nu}$	2
$\mathcal{O}_{HB} = ig'(D^\mu H)^\dagger(D^\nu H)B_{\mu\nu}$	2
$\mathcal{O}_T = \frac{1}{2}(H^\dagger \overset{\leftrightarrow}{D_\mu} H)^2$	1
$\mathcal{O}_{Hu} = (iH^\dagger \overset{\leftrightarrow}{D_\mu} H)(\bar{u}_R\gamma^\mu u_R)$	1
$\mathcal{O}_{Hd} = (iH^\dagger \overset{\leftrightarrow}{D_\mu} H)(\bar{d}_R\gamma^\mu d_R)$	1
$\mathcal{O}_{He} = (iH^\dagger \overset{\leftrightarrow}{D_\mu} H)(\bar{e}_R\gamma^\mu e_R)$	1
$\mathcal{O}_{HQ} = (iH^\dagger \overset{\leftrightarrow}{D_\mu} H)(\bar{Q}_L\gamma^\mu Q_L)$	1
$\mathcal{O}'_{HQ} = (iH^\dagger \sigma^a \overset{\leftrightarrow}{D_\mu} H)(\bar{Q}_L\sigma^a\gamma^\mu Q_L)$	1

with $F = \{L_L, e_R, Q_L, u_R, d_R\}$, Y_F the hypercharge, and

$$\mathcal{O}_{HL} \equiv (iH^\dagger \overset{\leftrightarrow}{D_\mu} H)(\bar{L}_L\gamma^\mu L_L),$$

$$\mathcal{O}'_{HL} \equiv (iH^\dagger \sigma^a \overset{\leftrightarrow}{D_\mu} H)(\bar{L}_L\sigma^a\gamma^\mu L_L). \tag{6}$$

Indeed, some important features of the dimension-six Lagrangian are best highlighted in other bases. In particular, the substitution $\mathcal{O}_{WW} \to \mathcal{O}_{HW}$ results in the strongly interacting light Higgs (SILH) basis of Refs. [10,19,22], which better captures the low-energy effects from universal UV theories (where the new physics only couples to SM bosons).[3] The substitution $\{\mathcal{O}_W, \mathcal{O}_B, \mathcal{O}_{HB}\} \to \{\mathcal{O}_{HL}, \mathcal{O}'_{HL}, \mathcal{O}_{WB}\}$ using Eqs. (3) and (5) leads to the basis of Ref. [7] (GIMR in what follows). IThe GIMR basis

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[3]In Ref. [19] also the operator $\mathcal{O}_r$ was replaced by $\mathcal{O}_H \equiv \partial_\mu|H|^2\partial^\mu|H|^2/2$ through a field redefinition: $2\mathcal{O}_r = \sum_{u,d,e}\mathcal{O}_y - 2\mathcal{O}_H + 4\mathcal{O}_6$.

mostly includes vertex corrections,[4] which makes the connection between operators and observables more straightforward [9]. Furthermore, nonuniversal theories in which new physics couples to the different fermions independently are more easily matched to this basis.

To understand which operators should be included in Higgs physics studies, we will now briefly discuss which ones are constrained by LEP using the basis of Table I (see Refs. [8,9] for detailed analyses in the bases of Refs. [7,19]).

The operators on the lhs of Table I are all of the form $|H|^2 \times \mathcal{L}_{SM}$. In the vacuum ($|H|^2 = v^2/2$), they merely redefine the SM input parameters and, thus, at tree level only contribute to Higgs physics [10]. All these operators modify the Higgs vertices and can be constrained by measuring the decay rates $h \to \gamma\gamma, Z\gamma, \bar{b}b, \bar{\tau}\tau$, the production modes $gg \to h, VV \to h, pp \to \bar{t}th$ and the trilinear h^3 coupling. At present, however, only the operators $\mathcal{O}_{BB}, \mathcal{O}_{WW}, \mathcal{O}_{GG}$ are constrained tightly enough to justify the EFT expansion (see Sec. III). On the other hand, the operators on the rhs of Table I also affect physics in the vacuum: $\mathcal{O}_{He}, \mathcal{O}_{Hu}, \mathcal{O}_{Hd}, \mathcal{O}_{HQ}, \mathcal{O}'_{HQ}$ and the combination $\mathcal{O}_W + \mathcal{O}_B$ are constrained by LEP measurements of Z-boson couplings to quark and leptons on the Z-pole (for the sake of counting, one can think of LEP1 as measuring independently the seven couplings of Z to ν, $e_{L,R}, u_{L,R}$ and $d_{L,R}$). These are all tightly constrained [8,23] and will be neglected in what follows. The operator $\mathcal{O}_{HB}$ and the combination $\mathcal{O}_W - \mathcal{O}_B$, on the other hand, affect in particular triple gauge couplings (TGCs) in addition to Higgs physics. Measurements of TGCs from diboson production at LEP2 and at the LHC constrain these operators. However, an analysis in the context of dimension-six operators including all existing data is not yet available (see the discussion in Ref. [27]). For this reason we include these operators when studying Higgs physics.

Hence, the Wilson coefficients of all operators in Table I are related to some experiment. Consequently, it is a prediction from $\mathcal{L}_6$ that any *additional* observable which can be extracted from Higgs physics is already constrained at some level of precision [8,9,25]. This is true, for example, for observables contained in channels with $V^* \to Vh$ associated production ($V \equiv W^\pm, Z$), to which we now direct our attention. Since the s-channel vector is off-shell, measurements of the differential distributions in these processes can access regions of momenta where the contribution of some operators is enhanced w.r.t. their contribution in gluon fusion and Higgs decays. Indeed, amplitudes such as $q\bar{q} \to V_L h$ involving longitudinal

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[4]Although the operator $\mathcal{O}_{WB}$ does contribute to the $W_3 B$ propagator, a combination of the operator $\mathcal{O}_{WB}$ and operators that modify gauge-boson/fermion vertices is unconstrained by LEP1 and is bound only at LEP2 because of its contribution to triple gauge vertices [9].
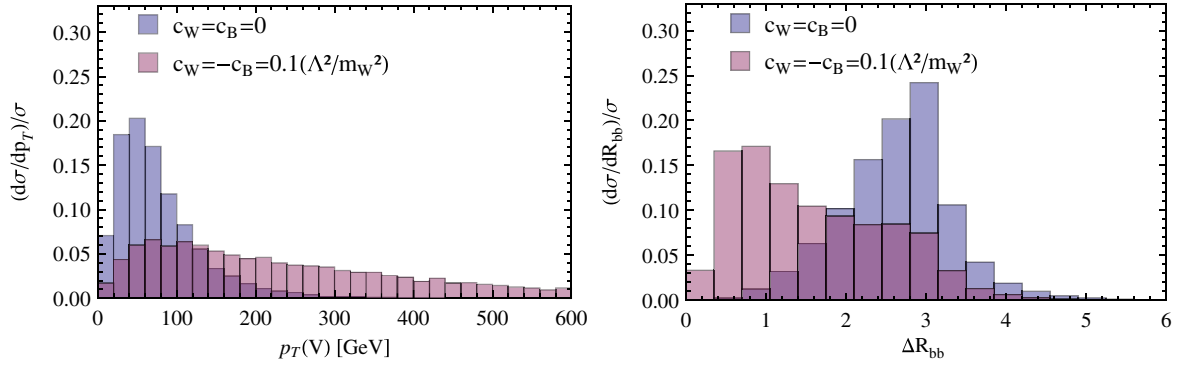
FIG. 1 (color online).   To illustrate the UV behavior of the operator $\mathcal{O}_W - \mathcal{O}_B$, these plots contrast the LO distributions of $p_T(V)$ and $\Delta R(b, \bar{b})$ $(pp \to ZH @ 8$ TeV$)$ for the SM and SM $+ \mathcal{O}$ with a large Wilson coefficient. No cutoffs are applied.
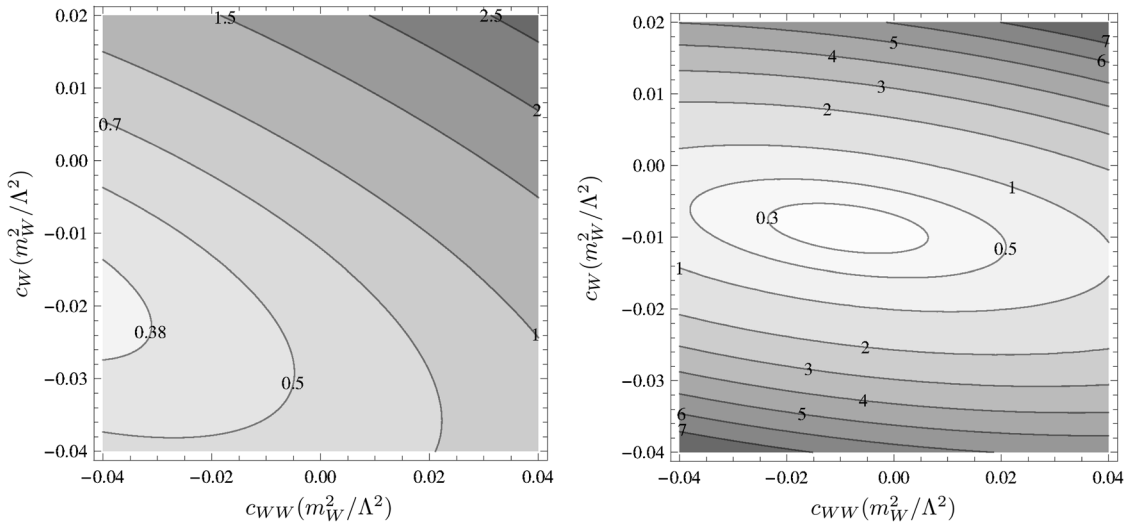

FIG. 2.   The impact of the operators $\mathcal{O}_{WW}$ and $\mathcal{O}_W$ on the cross section and kinematics of $pp \to Wh$ at the LHC8. Shown is $\sigma/\sigma_{\mathrm{SM}}$ (LEFT) and $\sigma/\sigma_{\mathrm{SM}}(p_T > 200)$ (RIGHT). The net effect of $\mathcal{O}_{WW}$ on the signal strength is subdominant in the region $p_T(W) > 200$ GeV. We assume that the EFT is valid up to the unitarity cutoff.

massive vector bosons will be sensitive to the breaking of gauge invariance communicated by the operators of Table I. In particular, unlike $\mathcal{O}_{VV}$, the operators $\mathcal{O}_V, \mathcal{O}_{HV}$ contribute to Goldstone boson production $q\bar{q} \to hG^{\pm,0}$ and $q\bar{q} \to G^{\pm,0}G^{\mp,0}$ and will, therefore, have the strongest impact on the high-energy tail of distributions in $V_L h$ and $V_L V_L$ final states. For $V = W^{\pm}$, only $\mathcal{O}_W$ and $\mathcal{O}_{WW}$ contribute to changes in kinematic distributions[5] as we can see by observing the squared partonic matrix element of $f\bar{f} \to W^+ h$.[6] In the SM, $|\mathcal{M}|^2 \to$ const for high energies at

tree level, and hence for $m_H \ll$ TeV, the amplitudes remain perturbatively unitary. However, in the presence of $\mathcal{O}_W$ and $\mathcal{O}_{WW}$, this changes and at large center-of-mass energy $\sqrt{\hat{s}}$,

$$\sum_T \int d\cos\theta |\mathcal{M}_T|^2 \to \frac{4g^4}{3}\frac{m_W^2}{\hat{s}}\left(1 + (c_{WW} + c_W)\frac{\hat{s}}{\Lambda^2}\right)^2,$$

$$\int d\cos\theta |\mathcal{M}_L|^2 \to \frac{g^4}{6}\left(1 + c_W\frac{\hat{s}}{\Lambda^2} + 4(c_{WW} + c_W)\frac{m_W^2}{\Lambda^2}\right)^2,$$

$$(7)$$

where we have separated the transverse polarizations of the W boson from the longitudinal one. As expected, the transverse ones are suppressed by a factor of $m_W^2/\hat{s}$, which is due to gauge invariance and the fact that the longitudinal polarization vector is proportional to p^μ/m_W in the high-energy limit. Note that the expansion makes sense only for

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[5]One linear combination of the other operators on the rhs of Table I can affect these processes through modifications of the Higgs branching ratios and wave-function normalization: we will comment on this in Sec. IV.

[6]The gluon-initiated process with the same final state might modify this cross section in the highly boosted regime [28].

$m_{W,h} \ll \sqrt{\hat{s}} \lesssim \Lambda$. From the formulae, we can also see that unpolarized measurements may mainly constrain the Wilson coefficient c_W because of its growing energy behavior in the linear part, while for $\mathcal{O}_{WW}$, due to its tranverse nature (field strength), the leading term is suppressed by m_W^2/Λ^2. We observe however that if we could single out the transverse polarization of the W boson, we could gain sensitivity on c_{WW} (see also [29]). Similarly, the operators $\mathcal{O}_B$, $\mathcal{O}_{HB}$ and $\mathcal{O}_{BB}$ will enter Z associated production.

The high-energy behavior of the cross section, described by Eq. (7), is portrayed in the lhs of Fig. 1, which shows the transverse momentum p_T distribution for pp collisions. The high-energy behavior also impacts the boost distribution of the Higgs in the laboratory reference frame, which is best captured by the $\Delta R(bb)$ distribution that we show in the rhs of Fig. 1. Notice from Eq. (7) that $\mathcal{O}_{BB}$ and $\mathcal{O}_{WW}$ give a smaller relative contribution to these processes at large $\hat{s}$, as we also illustrate in Fig. 2. For this reason, we concentrate our discussion on $\mathcal{O}_W$, $\mathcal{O}_B$ (actually, only the combination $\mathcal{O}_W - \mathcal{O}_B$ which is unconstrained by LEP1) and $\mathcal{O}_{HB}$ and comment later on generalizations.

III. ON THE VALIDITY OF THE EFT AT LARGE ENERGY

The EFT of Eq. (1) is an expansion in derivatives and SM fields over powers of Λ, defined as the scale where resonant new physics effects should become visible. Without additional assumptions, the EFT cannot be expected to describe processes at energies higher than Λ as operators of arbitrary dimension are then expected to become equally important, leading to a breakdown of the EFT description. In a bottom-up approach (from an IR point of view), Λ is not known a priori, but is a free parameter which needs to be fixed by experiment. The question whether or not the energy at which an experiment is performed lies within the validity of the EFT then depends on the sensitivity of the experiment itself. For instance, LEP1, working at c.o.m. energy $\sqrt{\hat{s}} = m_Z$, put bounds $\Lambda \gtrsim 1.6$ TeV for operators like the combination $\mathcal{O}_W + \mathcal{O}_B$. The sensitivity of the measurement hence fully justifies the EFT expansion in E/Λ, making the procedure self-consistent.[7] As we will see, at least for the Higgs production data available from the 7 and 8 TeV LHC runs, the situation is less clear.

Dimension-six operators including more derivatives with respect to an existing dimension-four interaction [class 2 in the classification of Eq. (2)] are expected to contribute an extra factor of $p^2 \sim \hat{s}$ to the amplitude compared to the SM and, hence,

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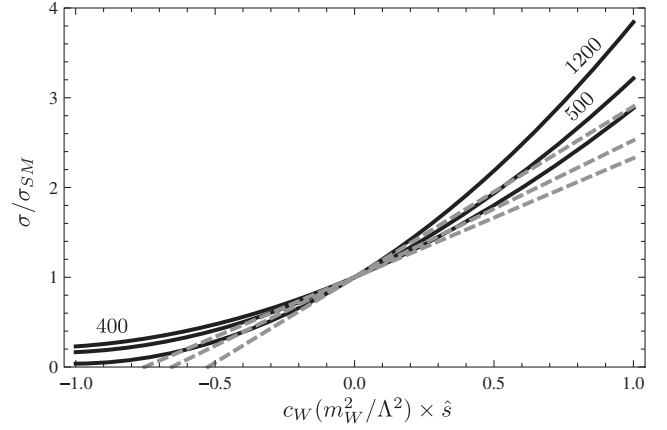[7]As we shall see below, the Wilson coefficient for these operators is typically $c \simeq 1$.

FIG. 3. The c_W dependence of the $u\bar{d} \to hW^+$ cross section at different fixed c.o.m. energies $\sqrt{\hat{s}} = 400, 500, 1200$ as described in the text. All orders of c_W are included in the squared amplitude for the solid lines. The dashed lines represent the linearized signal strengths.

$$\frac{\sigma}{\sigma_{\text{SM}}} \sim \left(1 + c_{i_2}\frac{\hat{s}}{\Lambda^2}\right)^2 \qquad (8)$$

(in reality, this somewhat simplistic view will be complicated by helicity effects). For $c_{i_2} \sim O(1)$, the points at which SM amplitudes are overtaken by EFT effects would typically mark the breakdown of the expansion in E/Λ. This is indeed the case for the operators in which we are interested. This is illustrated in Fig. 3, where we show the $u\bar{d} \to hW^+$ cross section in the presence of $\mathcal{O}_W$ at fixed center-of-mass energies $\sqrt{\hat{s}} = 400, 500, 1200$, and compare the first (linear) term of $\sigma/\sigma_{\text{SM}}$ in the $c_W E^2/\Lambda^2$ expansion with the complete expression. As expected, at c.o.m. energies $\sqrt{\hat{s}} \sim \Lambda/\sqrt{c_W}$, both the linear $O(c_W)$ and quadratic $O(c_W^2)$ contributions of the dimension-six Lagrangian to the cross section become comparable to the SM piece (for $c_W < 0$, the linearized signal strengths vanish already before this point, marking the lower limit of validity of this approximation). In this case, the question of the validity of the EFT is, therefore, related to the size of the EFT effects relative to the SM.

Despite the limitations of generic EFTs, most candidates for underlying models possess a more complicated structure in terms of different masses, couplings, and particle multiplicities. Hence, some of the Wilson coefficients of Eq. (2) might be parametrically larger (or smaller) for different operators. In particular, we have already mentioned in the previous section that under the assumption that new physics is characterized by a strong coupling g_*, the effective suppression of the operators of class 1 in Eq. (2), is

$$f \equiv \frac{\Lambda}{g_*} \qquad (9)$$

in the Lagrangian $\mathcal{L} \sim \mathcal{O}_1/f^2 + \cdots$, as these operators imply an expansion in fields which is valid only for small field values: $v/f \ll 1$. The important point is that f can be parametrically (up to $1 < g_* \lesssim 4\pi$ times) smaller than the masses of new particles $m_* \equiv \Lambda$, which mark the actual breakdown of the EFT. Cross sections which receive contributions growing with energy from these operators would make for a perfect probe for new physics. Indeed,

$$\frac{\sigma}{\sigma_{\text{SM}}} \sim \left(1 + c_{i_1} \frac{g_*^2(E)}{g_{\text{SM}}^2} \right)^2, \tag{10}$$

where g_{SM} describes the relevant (weak) SM coupling and we have defined $g_*(E) \equiv E/f$ [30]. From Eq. (9) we see that $g_*(E) < g_*$ for the EFT to be within the realm of validity but, contrary to Eq. (8), for $g_* \gg g_{\text{SM}}$ the EFT contribution relative to the SM can now be much bigger than unity, without exiting the realm of validity of the EFT. This situation arises, e.g., for 4-fermion operators and their contribution to $2 \to 2$ scattering growing with energy: under the assumption that g_* is large, it is possible to study this process at very high energy and obtain very tight constraints on the Wilson coefficients of these operators [31].

Do these arguments also apply to the operators that enter in Higgs physics? To answer this question in more detail, we study an explicit model with a spin one vector resonance V_μ^a, triplet under $SU(2)_L$, characterized by a coupling $g_* \lesssim 4\pi$ that might or might not be strong (similar arguments can be made with scalar or fermionic heavy states). Beside correctly describing all weakly coupled UV theories, our simplified model also captures the essence of strongly coupled scenarios that admit a weakly coupled holographic description (it is equivalent to a two-site model) in which V is a vector resonance emerging from the strong sector; the hope is that this description also qualitatively captures large classes of genuinely strongly coupled scenarios [19,32]. Heavy vectors have the correct quantum numbers to mediate interactions between Higgs, gauge boson or fermion currents, defined as

$$J_\mu^H \equiv \frac{i}{2} H^\dagger \overset{\leftrightarrow}{D}_\mu H, \qquad J_\mu^{Ha} \equiv \frac{i}{2} H^\dagger \sigma^a \overset{\leftrightarrow}{D}_\mu H,$$

$$J_\mu^{-H} \equiv \frac{i}{2} H^T \epsilon \overset{\leftrightarrow}{D}_\mu H, \tag{11}$$

$$J_\mu^a \equiv (D^\nu W_{\nu\mu})^a, \qquad J_\mu \equiv \partial^\nu B_{\nu\mu}, \qquad J_\mu^F \equiv \bar{F} \gamma_\mu F,$$

$$J_\mu^{Fa} \equiv \bar{F} \gamma_\mu \sigma^a F, \tag{12}$$

where F denotes any (chiral) fermion field. As sketched above, the most favorable scenario where these effects can be large is the one of strongly coupled theories, from which V emerges as a composite resonance. Then, any other composite state will couple to V with strength g_*. In particular, this is true for the Higgs field (a light composite

Higgs can arise as a pseudo-Goldstone boson from the strongly interacting sector [33]); SM gauge bosons, on the other hand, are very likely to be (mostly) elementary and in what follows we will assume that they form a separate (elementary) sector (the theoretical difficulties to realize composite gauge bosons are summarized in Ref. [34]). Then, the most general dimension-four Lagrangian[8] describing universal couplings of V yields

$$\mathcal{L}_{\text{Universal}} = \frac{m_*^4}{g_*^2} \left[\frac{1}{2m_*^2} V^{a\mu} V_\mu^a + \frac{V^{a\mu}}{m_*} \left(\gamma_H \frac{J_\mu^{Ha}}{m_*^3} + \gamma_V \frac{J_\mu^a}{m_*^3} \right) \right.$$
$$\left. - \frac{V^{a\mu\nu} V_{\mu\nu}^a}{4m_*^4} + \frac{|D_\mu H|^2}{m_*^4} \right], \tag{13}$$

where we have included the Higgs among the strongly coupled states and, in the spirit of naive dimensional analysis (NDA), have written the Lagrangian in a way that keeps the scaling in powers of masses and couplings manifest (see footnote 2). If the gauge fields, on the other hand, belong to a separate (elementary) sector, they are characterized as usual by $\mathcal{L}_{\text{el}} = -W_{\mu\nu}^a W^{a\mu\nu}/(4g^2)$ in non-canonical form and, similarly, for B_μ, G_μ^a.

In nonuniversal theories, the BSM sector can also couple to fermions. If the fermions (or combinations thereof) are also composite, we can write

$$\mathcal{L}_F = \sum_F \left(\gamma_F V^{a\mu} \frac{J_\mu^{Fa}}{g_*^2} + \frac{i}{g_*^2} \bar{F} \partial\!\!\!/ F \right). \tag{14}$$

If they are, however, elementary (or partially composite), the strong coupling g_* in front of the kinetic term in Eq. (14) should be replaced with the appropriate weak coupling. The coefficients $\gamma_{H,V,F} \sim O(1)$ quantify the departure from NDA, where they are expected to be of order unity. For canonically normalized fields, $V \to g_* V$, $H \to g_* H$, $F \to g_* F$ and $W \to gW$, we obtain the Lagrangian for V [35,36]:

$$\mathcal{L} = \frac{m_*^2}{2} V^{a\mu} V_\mu^a + V^{a\mu} \left(\gamma_H g_* J_\mu^{Ha} + \gamma_V \frac{g}{g_*} J_\mu^a + \sum_F \gamma_F g_* J_\mu^{Fa} \right)$$
$$- \frac{1}{4} V^{a\mu\nu} V_{\mu\nu}^a. \tag{15}$$

Then, integrating out the heavy vector triplets gives

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[8]Despite the appearance, the Lagrangian Eq. (13) can be associated to a renormalizable theory based on local gauge invariance, where V acquires its mass via a Higgs mechanism [35,36].

$$\mathcal{L} = -\frac{1}{2m_*^2}\left(\gamma_H g_* J_\mu^{Ha} + \frac{g}{g_*}\gamma_V J_\mu^a + \sum_F \gamma_F g_* J_\mu^{Fa}\right)^2 + \cdots$$

$$= c_r \frac{g_*^2}{m_*^2}\mathcal{O}_r + c_y \frac{g_*^2}{m_*^2}\mathcal{O}_y + c_6 \frac{g_*^2}{m_*^2}\mathcal{O}_6 + \sum_{F=Q,L} c'_{HF}\frac{g_*^2}{m_*^2}\mathcal{O}'_{HF} + c_W \frac{\mathcal{O}_W}{m_*^2} + c_{2W}\frac{\mathcal{O}_{2W}}{m_*^2} + \mathcal{O}_{4\text{fermi}} + \cdots, \quad (16)$$

where we expand in inverse powers of m_*, define

$$c_r = -\frac{3}{4}\gamma_H^2, \qquad c_y = \frac{\gamma_H^2}{8}, \qquad c_6 = \frac{\gamma_H^2}{2}, \qquad c'_{HF} = -\frac{1}{2}\gamma_H\gamma_F + \frac{1}{2}\gamma_V\gamma_F\frac{g^2}{g_*^2}, \qquad c_W = \gamma_H\gamma_V, \qquad c_{2W} = -\frac{\gamma_V^2}{2}\frac{g^2}{g_*^2}. \quad (17)$$

and introduce the operators $\mathcal{O}_{2B} \equiv (\partial_\mu B^{\mu\nu})^2$ and $\mathcal{O}_{2W} \equiv (D^\mu W_{\mu\nu}^a)^2$ and $\mathcal{O}_{4\text{fermi}}$.[9] The latter is denoting 4-fermion operators irrelevant for our discussion. The dots in Eqs. (16) denote higher derivative terms resulting from the momentum expansion in the propagator of V.

Similarly, we can study the effects of heavy vector singlets under $SU(2)_L$, but in order to avoid too large violations of custodial symmetry, we preserve the global $SU(2)_L \times SU(2)_R$ custodial symmetry of the SM and consider vectors $V^{0\mu}$, $V^{+\mu}$ triplets under $SU(2)_R$ [35]:

$$\mathcal{L} = \frac{m_*^2}{2}V^{0\mu}V_\mu^0 + m_*^2 V^{+\mu}V_\mu^- + V^{0\mu}\left(\delta_H g_* J_\mu^H + \delta_V \frac{g'}{g_*}J_\mu + \sum_F \delta_F g_* J_\mu^F\right)$$

$$+ \frac{1}{\sqrt{2}}\left(\delta_H g_* V^{+\mu}J_\mu^{-H} + \text{H.c.}\right) - \frac{1}{4}V^{0\mu\nu}V_{\mu\nu}^0 - \frac{1}{2}V^{+\mu\nu}V_{\mu\nu}^-. \quad (18)$$

In the low-energy theory, this yields the following coefficients:

$$c_r = -\frac{3}{4}\delta_H^2, \qquad c_y = \frac{\delta_H^2}{8}, \qquad c_6 = \frac{\delta_H^2}{2},$$

$$c_{HF} = -\frac{1}{2}\delta_H\delta_F + \frac{1}{2}\delta_V\delta_F\frac{g'^2}{g_*^2},$$

$$c_B = \delta_H\delta_V, \qquad c_{2B} = -\frac{\delta_V^2}{2}\frac{g'^2}{g_*^2}. \quad (19)$$

We are particularly interested in the coefficients of the operators $\mathcal{O}_W$, $\mathcal{O}_B$ and $\mathcal{O}_{HB}$. For the latter, it is clear that it does not arise at tree level from integrating out minimally coupled vectors, and its coefficient is, therefore, suppressed by a loop factor (similar arguments hold for $\mathcal{O}_{WW}$, $\mathcal{O}_{BB}$). One can, thus, estimate the coefficient suppressing the operator $\mathcal{O}_{HB}$ in the Lagrangian [19],

$$\mathcal{L}_{HB} \equiv \frac{c_{HB}}{\Lambda^2}\mathcal{O}_{HB}, \quad (20)$$

as

$$\frac{c_{HB}}{\Lambda^2} \simeq \frac{g_*^2}{16\pi^2 m_*^2} \lesssim \frac{1}{m_*^2}, \quad (21)$$

(up to factors of order one), where the inequality is saturated for maximally strongly coupled theories. Hence, these operators should not be trusted at energies higher than the inverse scale suppressing the operator; i.e., Eq. (20) should only be used at energies

$$E \lesssim \Lambda/\sqrt{c_{HB}} \simeq m_*, \quad (22)$$

and even then, this is true only for strongly coupled theories.

The operators $\mathcal{O}_W$, $\mathcal{O}_B$, on the other hand, do arise from vector exchange at tree level and (for elementary transverse gauge bosons) are not enhanced by a strong coupling: they are instead a genuine probe of the new physics resonance masses, as the coefficient that suppress them scales as

$$\frac{c_{W,B}}{\Lambda^2} \simeq \frac{1}{m_*^2}. \quad (23)$$

The effects they generate can be extrapolated only to energies $E \lesssim \Lambda/\sqrt{c_{W,B}}$, as was expected from Eq. (8). Naively, one might think that in the presence of a large number of nearly-degenerate vectors at the scale m_*, one could obtain an enhancement $c_{W,B} \sim N$ such that the effective scale that suppresses these operators could be $m_*/\sqrt{N} \ll m_*$ and, thus, parametrically smaller than the

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[9]The coefficients of $\mathcal{O}_{2B}$ and $\mathcal{O}_{2W}$ are proportional to $(g/g_*)^2$ or $(g'/g_*)^2$ since we have assumed the SM gauge bosons to be elementary. Had the gauge bosons been composite, then their contribution to $\mathcal{O}_{2B}$ and $\mathcal{O}_{2W}$ would have been tightly constrained by measurements of $e^+e^- \to e^+e^-$ at LEP2 as they correspond to the Y and W parameters [37].

cutoff. However, in minimally coupled UV scenarios where the heavy vectors are associated with additional spontaneously broken gauge symmetries, this is not the case. Indeed, in such a context, the coupling of the Higgs field to these vectors is characterized by its quantum numbers under these additional local symmetries: since the SM Higgs field only possesses four degrees of freedom, it cannot transform under N distinct $SU(2)$ symmetries, but only under a linear combination of them.[10] Furthermore, it is important to recall that the combination $\mathcal{O}_W + \mathcal{O}_B$, which for universal theories corresponds to the S parameter, is tightly constrained by LEP1 measurements. On top of this, in most interesting theories, the coefficients of these operators are strictly positive $c_{W,B} > 0$ [38], and consequently the combination $\mathcal{O}_W - \mathcal{O}_B$ which enters our analysis is already tightly constrained by LEP1.

Finally, for composite fermions, we see from Eqs. (16–19) that the operators $\mathcal{O}'_{HF}$ and $\mathcal{O}_{HF}$ are indeed enhanced by the strong coupling. In the Lagrangian,

$$\mathcal{L}_{HF} \equiv \sum g_*^2 \frac{c_{HF}}{\Lambda^2} \mathcal{O}_{HF} + g_*^2 \frac{c'_{HF}}{\Lambda^2} \mathcal{O}'_{HF}, \qquad (24)$$

the effective coefficient that multiplies each operator is

$$g_*^2 \frac{c_{HF}^{(\prime)}}{\Lambda^2} \simeq \frac{g_*^2}{m_*^2} = \frac{1}{f^2}, \qquad (25)$$

and the discussion of Eq. (10) applies: in particular, there exists a finite energy range $(g/g_*)m_* \lesssim E < m_*$ where the effect of these operators relative to the SM can be much bigger than one while the EFT expansion is still valid. That it is indeed still valid can be seen by looking at the form of dimension-eight operators that arise from Eq. (16): operators with more derivatives (schematically of the form $(p^2/m_*^2) \times \mathcal{O}_{HF}$ in momentum space) which contribute to the same tree-level process, originate from Eq. (16) at the next order in the momentum expansion, and their contribution to Eq. (10) is $\sim g_*^2 E^4/(g_{SM}^2 m_*^4) = g_*^2(E)/g_{SM}^2(E^2/m_*^2) = g_*^2(E)/g_{SM}^2(g_*(E)^2/g_*^2)$. This shows that the cutoff is indeed m_* and not $m_* g_{SM}/g_*$. Notice that this also implies that it is consistent to keep contributions of order $(c_{HF}^{(\prime)})^2$, since these are expected to be much bigger than the contributions from dimension-eight operators to the same process.

We have, thus, found a set of operators which can also be studied in a regime where their relative contribution to the SM amplitudes is much bigger than one. How do these operators contribute to Vh associated production? As discussed in the previous section, most of the operators

$\mathcal{O}_{HF}^{(\prime)}$ are already tightly constrained by LEP1 as they modify the couplings of the gauge bosons to fermions. Nevertheless, there is one combination of these operators which is equivalent to an overall shift of the Weinberg angle in the gauge-fermion sector and, as such, cannot be constrained by LEP1 [9]; it can only be measured as a relative shift between θ_W as measured in $Z\bar{F}F$ couplings and θ_W as measured in gauge bosons self-couplings or in Higgs physics. Indeed, this direction is equivalent to [9],[11]

$$\Delta\mathcal{L}_{F_{tot}} = 2\tan^2\theta_W\left(-\mathcal{O}_T + \sum_F Y_F \mathcal{O}_{HF}\right) - \mathcal{O}'_{HL} - \mathcal{O}'_{HQ}, \qquad (26)$$

which, using Eq. (5), can be shown to induce the same effects as [9,25],

$$\Delta\mathcal{L}_{F_{tot}} = \frac{4}{g^2}(\mathcal{O}_B - \mathcal{O}_W) + \mathcal{O}_{\theta_W(Higgs)}, \qquad (27)$$

where

$$\mathcal{O}_{\theta_W(Higgs)} = 6\mathcal{O}_r - \sum_{u,d,e}\mathcal{O}_y - 4\mathcal{O}_6 \qquad (28)$$

modifies the Higgs vertices independently of momentum. Indeed it can be easily seen that Eq. (27) contributes only to TGCs (in particular to the parameter g_Z^1 [39]) or Higgs physics. Interestingly, from the arguments given above, the contribution to $\mathcal{O}_W - \mathcal{O}_B$ from the particular direction Eq. (26) is enhanced by a g_*^2/g^2 factor with respect to the naive contribution from universal theories and provides a motivated context in which the effect of these operators can be studied at high energy, as discussed in Eq. (10).

This discussion of the breakdown of the EFT from a top-down perspective can be complemented with a bottom-up approach (without detailed knowledge of the UV theory) by analyzing perturbative partial wave unitarity. An analysis of partial wave unitarity violation for a several dimension-six operators has been performed in [40]. The operators $\mathcal{O}_{HW}$ and $\mathcal{O}_{HB}$ imply the constraints

$$\hat{s} \lesssim 15.5\frac{\Lambda^2}{c_{HW}}, \qquad 49\frac{\Lambda^2}{c_{HB}}. \qquad (29)$$

Since $\mathcal{O}_{HW}$ yields by far the strongest unitarity constraint, we use it as an estimate for the unitarity violation induced by $\mathcal{O}_W = \mathcal{O}_{HW} + \mathcal{O}_{WW} + \mathcal{O}_{WB}$. While universal EFTs are far away from saturating this bound, in the case of Eq. (26)

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[10]In fact, the four d.o.f. of the SM Higgs doublet can be cast into a $(\mathbf{2},\mathbf{2})$ of $SU(2)_{SM} \times SU(2)_{BSM}$ and can transform at most under one additional $SU(2)_{BSM}$ gauge group; this is the model of Eq. (15) [35,36].

[11]This direction involves a large contribution to the operator $\mathcal{O}_T$ which is generated by interactions with heavy fields that do not preserve custodial symmetry as in Eq. (18); nevertheless, the combination Eq. (26) is accidentally custodial preserving, as the custodial breaking part cancels between $\mathcal{O}_T$ and $\mathcal{O}_{HF}$.

we will restrict ourselves to values of g_* which satisfy these constraints.

In summary, we have found that for generic EFTs, extrapolation of the effects of the dimension-six operators in a regime where their contribution, relative to the SM, is bigger than one, is inconsistent with the EFT expansion itself. For the operators that can contribute to HV associated production, this is true also in universal theories characterized by a strongly coupled Higgs sector, where $\mathcal{O}_{HB}$ arises only at loop level, while $\mathcal{O}_W$ and $\mathcal{O}_B$ are suppressed by the cutoff itself; this situation does not improve in theories with many vectors. Nevertheless, in theories in which the particular combination of fermions reported in Eq. (26) is composite and part of a strongly coupled sector, the combination $\mathcal{O}_W - \mathcal{O}_B$ can be enhanced by the strong coupling, and its effects can be studied also in a regime where its relative contribution is much bigger than the SM one.[12]

IV. BOUNDS FROM EXISTING LHC HIGGS SEARCHES

As explained in Sec. II, Higgs associated production channels can probe Higgs interactions at high energy and are particularly sensitive to BSM interactions like $\mathcal{O}_B$, $\mathcal{O}_{HB}$ etc., whose contribution strongly increases with the center-of-mass energy [see Eq. (7)]. However, in the previous section we have shown that in generic EFTs, the perturbative expansion breaks down at large energy when the relative contribution of these operators is bigger than one. In this case, experiments whose sensitivity is of the order of the SM contribution or weaker will not be able to put meaningful constraints on the EFT. We have shown that the same arguments hold in universal theories even in the strong coupling limit. Within the relatively general framework which we have considered, namely that of perturbative minimally coupled UV completions, only specific scenarios with strongly interacting fermions allow us to to extrapolate the validity of the EFT at large energies, and for the operator $\mathcal{O}_W - \mathcal{O}_B$ only.

For this reason, we begin with a study of the $\mathcal{O}_W - \mathcal{O}_B$ direction. It is important to notice that a study of this combination in isolation (i.e. by assuming that the coefficients of all other operators are much smaller) makes sense for a number of reasons. First of all, $\mathcal{O}_W - \mathcal{O}_B$ is orthogonal to physics from LEP1, meaning that we can ignore LEP1 constraints in our discussion as well as the other operators that contribute to LEP1 observables: $\{\mathcal{O}_T, \mathcal{O}_{Hu}, \mathcal{O}_{Hd}, \mathcal{O}_{He}, \mathcal{O}_{HQ}, \mathcal{O}'_{HQ}\}$ and the combination $\mathcal{O}_W + \mathcal{O}_B$. Also, $\mathcal{O}_W - \mathcal{O}_B$ does not contribute to the $h \to \gamma\gamma$ or $h \to Z\gamma$ partial widths or the hgg coupling, all of which are tightly constrained from LHC measurements. Thus, beside the analysis of this article, the only constraints on

$\mathcal{O}_W - \mathcal{O}_B$ are from TGC measurements. Furthermore, from a theoretical point of view, we have argued that within the class of UV physics we consider, c_{HB}, c_{WW}, c_{BB} are very small despite the enhancement of $c_W - c_B$ and, in case their size would be big enough to be relevant for the experiment, then the EFT would not be valid. Finally, the contribution of other operators (such as $\mathcal{O}_r$ or all remaining operators in Table I that modify the Higgs width) does not grow as fast with energy in the Vh associated production cross section and their impact is negligible if their coefficients are within the validity of the EFT expansion [see below Eq. (9)].[13]

We will, therefore, study constraints on this combination of operators first and then discuss possible extension. We will repeat the analysis for different choices of UV realizations: (i) universal theories (or generic EFTs) where the scale that suppresses the operator is the cutoff[14] (i.e. Wilson coefficients of order unity) and (ii) theories with composite fermions, in which the Wilson coefficient can be large, implying a large hierarchy between the scale that suppresses the operator and the cutoff.

A naive comparison of the total cross section with measured signal strengths is inadequate: on the one hand, the effects of $\mathcal{O}_W - \mathcal{O}_B$ are strongest in high p_T bins which have the lowest SM+Higgs background, while on the other hand it is precisely those bins which might be probing the breakdown of the EFT. For this reason, the full differential distribution must be considered; we do so in this section and discuss the dependence of the bounds obtained on the choice of cutoff, which we take consistently into account. Indeed, for scenario (ii), we rely on the strong coupling and use data from energies up to the unitarity bound Eq. (29). For scenario (i), on the other hand, we must discard information coming from events whose energy lies beyond the region of generic EFT validity for any given value of the coefficients $(c_W - c_B)/\Lambda^2$. Since this affects almost exclusively kinematic regions where there is very little SM background, this cutoff reduces χ^2 and, thus, yields conservative exclusions.[15]

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[12]This is true also in theories where the gauge bosons are fully composite, but we have mentioned above the theoretical limitations and the experimental constraints of these theories.

[13]At present, the bounds on these operators from analyses of Higgs data are not strong enough to justify the EFT expansion (see Refs. [8,24,41,42]). As a consequence, also effects of order, e.g., $c_r c_W$ in the cross section for Vh, channels can be important: another sign of the breakdown of the perturbative expansion as these are clearly of the order dimension eight.

[14]While p_T and $\hat{s}$ distributions are related, events of any given p_T can in principle have arbitrarily high energies at large rapidities. Therefore, we will use $\hat{s}$ cuts to impose cutoffs in what follows. In practice, cutting on $\sum(m_i^2 + p_{iT}^2)^{1/2}$ is a reasonable approximation.

[15]An estimate of the uncertainty in results due to the breakdown of perturbativity can also be obtained by comparing the constraints obtained from linearized signal strengths $\sigma/\sigma_{\rm SM} \approx 1 + a c_W$ with the full result, which includes contribution of the same order as those of dimension-eight operators. Using a variable cutoff procedure as we do in this paper both yields more conservative results for the generic EFT case and allows the treatment of EFTs with enhanced Wilson coefficients in which $O(c^2)$ effects are actually physically meaningful.

We extract bounds on the coefficients of these operators using present data on Higgs associated production and concentrate on the final state with two b jets, leptons and missing energy [43,44]:

$$pp \rightarrow Zh; \qquad h \rightarrow b\bar{b}, \qquad Z \rightarrow l\bar{l}, \nu\bar{\nu}$$
$$pp \rightarrow W^{\pm}h; \qquad h \rightarrow b\bar{b}, \qquad W^{\pm} \rightarrow \bar{l}\nu/l\bar{\nu}. \quad (30)$$

We have implemented the corresponding ATLAS searches [43], where data and expected background and signal events for each p_T bin are reported. A caveat concerning the interpretation of the data given in Table V of [43] is in order. While a combined χ^2 fit of the total signal strength using the provided $0l, 1l, 2l$ data for observed, expected, background and background errors yields a reasonable confidence interval in line with the official combination, the best-fit values for individual channels differ from the official values. This might be due to correlations, in which case the necessary information to improve on this is currently unavailable to the public. Our combined results, which are dominated by the highest p_T bins below the cutoff, are expected to be less sensitive to this problem. The simulations are performed using MADGRAPH 5[45], PYTHIA [46], and DELPHES [47] using our FEYNRULES [48,49] implementation of the effective theory and the *cteq6l1* [50] PDF sets with variable factorization scale corresponding to the MG5 standard setting. The analyses of Ref. [43] use 5 (2*l* and 1*l*) or 3 (0*l*) different $p_T(V)$ bins separated at $p_T(V) = (0–90, 90–120), 120–160, 160–200, > 200$ GeV which are subject to different additional kinematic cuts. By treating these bins separately, we gain sensitivity to the shape of the p_T distributions, and in particular to the high-energy behavior of the EFT.

For generic EFTs [case (i)], $c_{W,B} \sim 1$ and the appropriate cutoff is the inverse of the scale suppressing the dimension-six operator. This means that to every value of the coefficient $1/\Lambda^2$ corresponds a different cutoff $E < \Lambda$. The larger the value of the coefficient, the smaller the amount of data available to constrain it. As illustrated by the blue curve in Fig. 4, where we plot the $\Delta\chi^2$ contour for $(c_W - c_B)/\Lambda^2$, the present sensitivity is only enough to put a constraint on these operators due to an underfluctuation in the data, while there is no expected limit as the corresponding $\Delta\chi^2$ never even passes the 2σ threshold. In the region $c_W = -c_B < 0$, neither the observed nor the expected $\Delta\chi^2$ yield exclusions. The same is true for the $\mathcal{O}_{HB}$ operator with the cutoff suggested by Eq. (21).

In case (ii), the validity of the EFT is extended to energies parametrically larger than the inverse of the scale suppressing the dimension-six operator. The $\Delta\chi^2$ contours obtained with this cutoff correspond to the purple, yellow and green curves in Fig. 4. As the discussion in the previous sections indicates, this parametric enhancement depends on the size of the strong coupling. Since the naive choice for a maximally strong coupling, $g_* = 4\pi$, violates the partial wave unitarity bound, we limit ourselves to values of g_* that respect Eq. (29) (notice, however, that for values of g_* as large as 4π the bounds do not change noticeably, see Fig. 4). The corresponding $\Delta\chi^2$ contours are given by the green curve in Fig. 4. We obtain the following consistent constraint on these operators:

$$-0.06 \lesssim \frac{c_W - c_B}{\Lambda^2/m_W^2} \lesssim 0.02, \qquad 95\% \text{ C.L.} \quad (31)$$
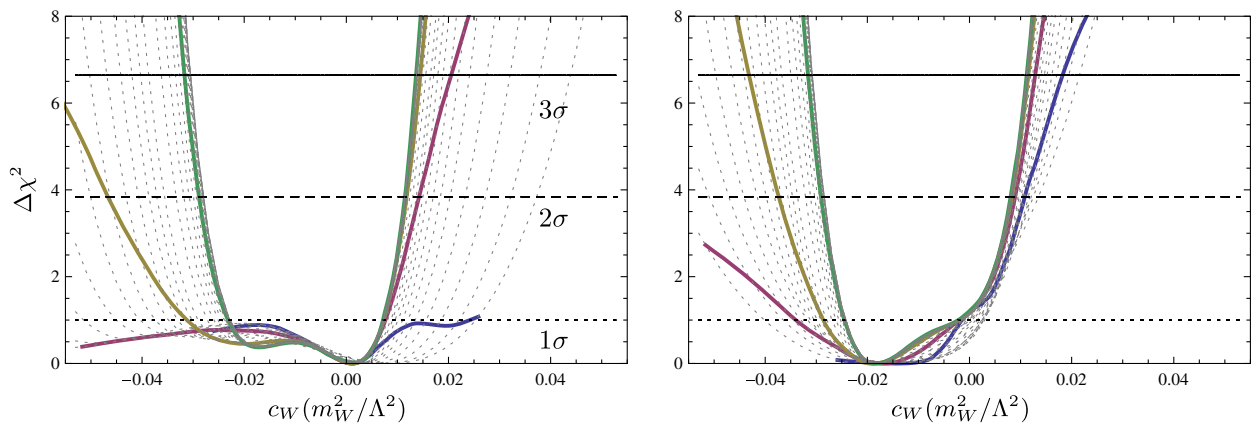

FIG. 4 (color online). The combined expected (left) and observed (right) 1-parameter fit $\Delta\chi^2$ contours in the coefficient $c_W(m_W^2/\Lambda^2) = -c_B(m_W^2/\Lambda^2)$ from Higgs searches in the $b\bar{b} + 0l, 1l, 2l$ final states in ATLAS. We assume all other operators in the basis to be negligible and employ various UV cutoff prescriptions. The dashed contours are for fixed UV cuts $\sqrt{\hat{s}} < 500, 550, \ldots$ GeV, while the solid contours are for parameter-dependent cutoffs $\hat{s} < \Lambda^2/c_W$ (blue, lowest lying curve), $2\Lambda^2/c_W$ (purple, next-to lowest curve), $4\Lambda^2/c_W$ (yellow, next-to-steepest curve) and $4\pi\Lambda^2/c_W$ (green, steepest curve) inspired by our discussion of UV completions and perturbativity. We assume that the main source of error is statistical and treat the theoretical errors as nuisances.

A. Comparison with TGCs

As explained in Sec. II, we want to quantify the added information of studying channels that probe Higgs physics at high energy. For this reason we have neglected all operators that are tightly constrained by either LEP1 or by measurements of $h \to \gamma\gamma, Z\gamma$ and $gg \to h$: then, only the combinations[16] [9,25]

$$\mathcal{O}_W - \mathcal{O}_B, \qquad \mathcal{O}_{\overline{HB}} \equiv \mathcal{O}_{HB} - \frac{1}{2}(\mathcal{O}_{WW} - \mathcal{O}_{BB}) \quad (32)$$

yield contributions to Higgs observables which grow fast with energy and are not tightly constrained by other experiments (since $\mathcal{O}_{HB}$ contributes to $h \to Z\gamma$, in Eq. (32) we have cancelled this contribution by subtracting a piece $(\mathcal{O}_{WW} - \mathcal{O}_{BB})/2$ that contributes to $h \to Z\gamma$ only [9,25]). As a matter of fact, these operators also modify TGCs (measured in $e^+ e^- \to W^+ W^-$ scattering at LEP2), so that it is tempting to compare which experiment gives the strongest constraints (see also Refs. [41,51]).

The sensitivity at LEP2 was high enough to constrain the Wilson coefficients in TGC measurements within the realm of perturbativity of generic EFTs. The EFT description at LEP2 is, therefore, adequately self-consistent. Nevertheless, there exists at present no analysis of LEP2 data which consistently includes the effects of all dimension-six operators (see the discussion in Ref. [27]). A sensible assumption which allows us to derive bounds on $\mathcal{O}_W - \mathcal{O}_B$ and $\mathcal{O}_{HB}$ from TGC measurements is to limit ourselves to a generic class of theories where the operator $\mathcal{O}_{3W} = \frac{\epsilon^{abc}}{3!} W_\mu^{a\nu} W_{\nu\rho}^b W^{c\rho\mu}$ is small. Under this assumption, the 95% C.L. bounds from TGCs are [52][17]

$$-0.05 \lesssim \left(\frac{c_W - c_B}{2}\right)\frac{m_W^2}{\Lambda^2} \lesssim 0.05,$$

$$-0.12 \lesssim c_{\overline{HB}}\frac{m_W^2}{\Lambda^2} \lesssim 0.10. \quad (33)$$

Note that this upper bound on c_W from LEP corresponds to a suppression scale $\gtrsim 350$ GeV, larger than relevant LEP2 energies.

On the other hand, as discussed above, the constraints from Higgs observables at high energy that we have derived here are typically beyond the validity of the EFT expansion, but they can make sense for the direction $\mathcal{O}_W - \mathcal{O}_B$, in the

FIG. 5 (color online).   The 95% C.L. (solid) and 99% C.L. (dashed) combined observed limits on the coefficients c_W and $c_{\bar{H}B}$ (with $c_B = -c_W$ and all other operators set to zero) from our analysis of Higgs searches in the $b\bar{b} + 0l, 1l, 2l$ final states in ATLAS. We employ a cut $\sqrt{\hat{s}} < 1.2$ TeV. We compare the exclusion with LEP2 limits on TGCs (red contour).

case of strongly interacting fermions. Nonminimally coupled theories could in principle generate tree-level effects for c_{HB}, but it is difficult to argue along the lines of Sec. III to say whether the coefficient of these operators can or cannot be enhanced with respect to the inverse cutoff. We assume for completeness that a class of theories exists where the coefficients of the operator $\mathcal{O}_{HB}$ can be very large, and that the validity of the EFT description can be extrapolated up to the breakdown of perturbative unitarity. The resulting bounds from present Higgs data, valid only in this class of theories, are shown in Fig. 5. We employ a cut $\sqrt{\hat{s}} < 1200$ GeV corresponding to $\sqrt{4\pi}m_W/\sqrt{0.05}$, keeping however in mind that unlike for $c_W - c_B$, large values of c_{HB} as they appear in this fit are not described in terms of the UV models presented before. As mentioned above, we are showing the direction $\mathcal{O}_{\bar{H}B}$ rather than $\mathcal{O}_{HB}$ only, as the former gives results that are independent from bounds on $h \to Z\gamma$ [9,25]: in this way the two-dimensional plot shown in Fig. 5 is a genuine comparison between TGCs and Higgs physics at high energy, and is unaffected by bounds from any other experiment at present.[18] For this reason, we believe that the plot of Fig. 5 is particularly instructive; furthermore, it quickly allows us to differentiate (between vertical and horizontal axes) along which direction the comparison with

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[16]Due to a different convention in the covariant derivative, our definition of $\mathcal{O}_{HB}$ differs from Ref. [25].

[17]In our basis, the TGC parameters of Ref. [39] are modified as $\delta g_1^Z = (c_W/\cos^2\theta_W)(m_W^2/\Lambda^2)$ and $\delta\kappa_\gamma = c_{HB}(m_W^2/\Lambda^2)$. As noticed in Ref. [27], under the assumption that $c_{3W} = 0$, there is no quantitative difference between a fit to TGCs in the context of dimension-six operators (that neglects terms higher order in the Wilson coefficients) and the fit of the LEP2 collaboration [52], which we use in this article.
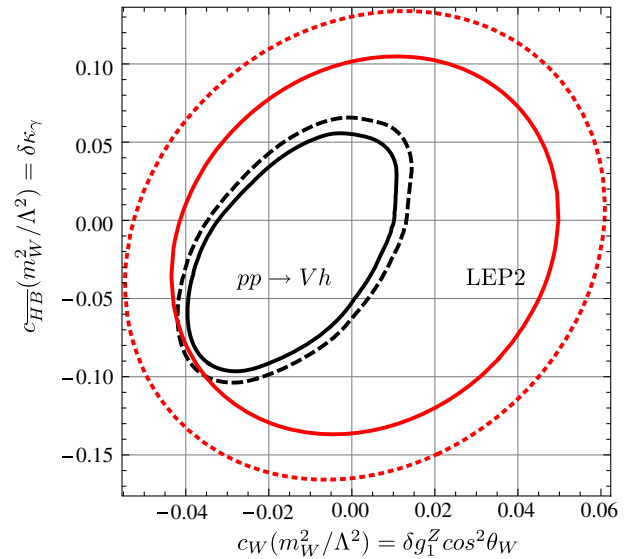
[18]Due to the large coefficients in front of c_{WW} in Eq. (7), the part $(\mathcal{O}_{WW} - \mathcal{O}_{BB})$ has, nevertheless, a sizable impact on the HV channel, although it doesn't grow fast with energy; for this reason, we differentiate between $\mathcal{O}_{HB}$ and $\mathcal{O}_{\bar{H}B}$.

TGC makes sense (and along which one it does not) within the context of the theories described above.

V. CONCLUSIONS

We have investigated constraints on new physics from LHC Higgs searches in an EFT context, with an emphasis on channels that are sensitive to high energies (in particular Higgs associated production $pp \to hV$) and can potentially be very good probes to search for new physics at hadron machines. In such Higgsstrahlung processes, the invariant mass is only limited by PDF suppression (as opposed to on-shell Higgs production, where $\hat{s} \approx m_h^2$) and cross sections as well as distributions can be drastically modified by the presence of dimension-six operators in the Lagrangian. Indeed, some operators, unconstrained by LEP1 and by measurements of on-shell Higgs properties, and only mildly constrained by LEP2 TGC measurements, contribute to the effective hVV vertex in a manner that grows with energy.

We have concentrated on these operators (namely, the combinations $\mathcal{O}_W - \mathcal{O}_B$ and $\mathcal{O}_{\bar{H}B}$ in the basis of Table I) and discussed the extent to which EFT analyses of LHC Higgs searches can sensibly use the high-energy tail of distributions to exploit this growth. In particular, we have shown that in the context of universal theories (where new physics couples—strongly or weakly—to the SM bosons only) as well as in theories characterized only by a scale Λ and weak couplings, the EFT expansion is not valid at such large energy. If a consistent analysis, suitable for universal theories, is performed, then the present data is not accurate enough to provide any constraints on dimension-six operators.

The very essence of using EFTs to parametrize and constrain new physics BSM is that they can describe large classes of UV scenarios in a simple way and allow us to quickly reinterpret bounds on the Wilson coefficients as bounds on masses and couplings of BSM particles. For this reason, rather than simply assuming the existence of UV scenarios for which the EFT expansion is valid also at high energy, it is crucial to understand if these scenarios really exist and which assumptions they require. Therefore, we have explicitly constructed a class of UV models, characterized by a strong coupling g_* in addition to the scale Λ, to study under which circumstances and for which operators the reach of EFTs can be extended. We have found that within this relatively general class of models, one particular combination of SM fermions needs to be strongly coupled (e.g. as composites emerging from a strongly coupled sector). This scenario would have been impossible to constrain at LEP1, but can be constrained by TGC measurements or through the analysis we present here. This is the only concrete scenario we have found in which a study of the differential distribution of the $pp \to hV$ channels (and in particular of their high-energy tail) can provide strong and consistent constraints on Wilson coefficients. These constraints are complementary to LEP1 in the context of fermion compositeness, and are competitive with LEP2.

Furthermore, the indirect limits on anomalous TGCs derived from Higgsstrahlung using various cutoff prescriptions can serve as a consistency check between direct searches for new physics and anomalous TGC measurements.

Finally, the searches outlined in this paper can play an important role in future high-energy and high-luminosity runs of the LHC, where more precise measurements of Higgs and gauge boson production rates and kinematics will compete with direct searches to constrain or discover new physics (see also [21]).

Note added.—Recently, Ref. [20] appeared which has some overlap with the present work. While numerical results agree with the v3 of Ref. [20] where comparable, our detailed analysis of the EFT breakdown, its impact on LHC Higgs searches and the interpretation in terms of UV completions are unique to our work.

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