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Towards the ultimate SM fit to close in on Higgs physics

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ABSTRACT: With the discovery of the Higgs at the LHC, experiments have finally addressed all aspects of the Standard Model (SM). At this stage, it is important to understand which windows for beyond the SM (BSM) physics are still open, and which are instead tightly closed. We address this question by parametrizing BSM effects with dimension-six operators and performing a global fit to the SM. We separate operators into different groups constrained at different levels, and provide independent bounds on their Wilson coefficients taking into account only the relevant experiments. Our analysis allows to assert in a model-independent way where BSM effects can appear in Higgs physics. In particular, we show that deviations from the SM in the differential distributions of $h \rightarrow V\bar{f}f$ are related to other observables, such as triple gauge-boson couplings, and are then already constrained by present data. On the contrary, $BR(h \rightarrow Z\gamma)$ can still hide large deviations from the SM.

KEYWORDS: Higgs Physics, LEP HERA and SLC Physics, Beyond Standard Model

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1 Motivation

With the first LHC revenues we have begun to gain experimental evidence for all sectors of the Standard Model (SM). Not having found signs of beyond the SM (BSM) physics, it is time to ask which directions in the parameter space of “deformations” of the SM are still unprobed. This can provide a useful guidance for future experiments.

The purpose of this paper is to take steps towards addressing this question. We will mainly focus on flavor-independent observables, and how these constrain electroweak symmetry breaking (EWSB) effects in gauge-bosons and Higgs physics. As a model-independent parametrization of BSM physics, we will use the Wilson coefficients of the independent dimension-six operators. This is a valid parametrization as long as the BSM scale is heavier than the electroweak scale, as we will assume here. Instead of an exhaustive fit to the SM, including all available data, our goal is to show which are the most relevant experiments that constrain the different directions in the parameter space of Wilson coefficients, and provide the corresponding bounds.

We will classify the bounds in different groups [1]. First, those arising from Z -pole observables, W mass and some low-energy experiments that constrain, at the per-mille level, deviations in W/Z propagators and gauge-boson couplings to fermions. In a second group, we have bounds, a factor ~ 10 weaker, from measurements of the triple gauge-boson couplings (TGC), that will be crucial to constrain certain directions in the parameter space,

left unbound by the first group of experiments. Finally, we will have bounds from Higgs physics arising from the recent LHC data.

One of the main purposes of our fit is to extract indirect constraints on possible deviations from the SM in future Higgs measurements. In particular, we will find that only $\Gamma(h \rightarrow Z\gamma)$ can still hide large deviations from the SM prediction, while deviations in $h \rightarrow V\bar{f}f$ are already constrained by other experiments, as they can be related to TGC measurements. Since the LHC can in principle do better measurements of TGC from processes such as $pp \rightarrow W\gamma, Z\gamma, WW, ZZ$, this will be a more optimal way to constrain BSM physics than from $h \rightarrow V\bar{f}f$ or $VV \rightarrow h$.

There have been several groups also addressing these questions [2–8]. Contrary to ours, however, most analyses have considered operators one by one, setting the rest artificially to zero, and the combined effect of more than one operator is rarely taken into account. The analysis that goes closer in spirit to the one presented here is the one of ref. [2]. Our analysis differs from theirs in several aspects. First of all, our analysis is extended to include Higgs physics and tries to understand where BSM effects can be found in future Higgs measurements. Secondly, and importantly, our choice of basis [1, 9, 10] allows to separate operators that are constrained at different level. This choice minimizes the correlations between Wilson-coefficient constraints and makes the results easy to interpret. The analysis of ref. [2] was made using the basis of ref. [11], and led to ref. [7] to erroneously conclude that constraints from electroweak observables were much weaker than what they actually are.

In section 2 we study the constraints on the CP-even Wilson coefficients arising from electroweak (non-Higgs) observables. In section 3 we present the relevant Wilson coefficients for Higgs physics, show their constraints, and study new physics in $h \rightarrow V\bar{f}f$. In section 4 we briefly comment on CP-odd operators. We leave for appendix A the details concerning Z -pole observables, and in appendix B we show how LHC data can constrain the relevant four-fermi interactions entering in our analysis.

2 Constraining BSM physics from EW observables

We will assume that the BSM sector is heavy, with a mass scale Λ much larger than the weak-scale. This allows us to parametrize in a model-independent way all BSM effects by dimension-six operators added to the SM [11, 12]. There are many options for the choice of the dimension-6 operator basis. Here we will follow the basis of [1, 10] that we find more suitable for our analysis than that of ref. [11], due to its better connection with experiments. Assuming lepton and baryon number conservation, the full set of operators in this basis can be found in tables 1 and 2 of ref. [1]: these contain 5 redundancies that we use to eliminate the 5 operators $\mathcal{O}_{2W,2B,2G}$ and $\mathcal{O}_L^i, \mathcal{O}_L^{(3)i}$.

In table 1 we show the dimension-6 operators of our basis that enter in the observables that we are interested in. We assume minimal flavor violation (MFV) [13], that is to say that the BSM sector respects a $U(3)^5$ flavor-symmetry up to corrections proportional to the SM Yukawa couplings. This allows us to concentrate on flavor-independent constraints. In particular, MFV implies that dipole operator for fermions, as well as $(i\tilde{H}^\dagger \overleftrightarrow{D}_\mu H)(\bar{u}_R \gamma^\mu d_R)$, are proportional to small Yukawa couplings and can then be neglected. In the operators

$\mathcal{O}_H = \frac{1}{2}(\partial^\mu H ^2)^2$ $\mathcal{O}_T = \frac{1}{2}(H^\dagger \overleftrightarrow{D}_\mu H)^2$ $\mathcal{O}_6 = \lambda H ^6$	$\mathcal{O}_{BB} = g'^2 H ^2 B_{\mu\nu} B^{\mu\nu}$ $\mathcal{O}_{GG} = g_s^2 H ^2 G_{\mu\nu}^A G^{A\mu\nu}$ $\mathcal{O}_{HW} = ig(D^\mu H)^\dagger \sigma^a (D^\nu H) W_{\mu\nu}^a$ $\mathcal{O}_{HB} = ig'(D^\mu H)^\dagger (D^\nu H) B_{\mu\nu}$ $\mathcal{O}_{3W} = \frac{1}{3!} g \epsilon_{abc} W_\mu^{a\nu} W_{\nu\rho}^b W^{\rho\mu c}$
$\mathcal{O}_W = \frac{ig}{2} (H^\dagger \sigma^a \overleftrightarrow{D}_\mu H) D^\nu W_{\mu\nu}^a$ $\mathcal{O}_B = \frac{ig'}{2} (H^\dagger \overleftrightarrow{D}_\mu H) \partial^\nu B_{\mu\nu}$	

$\mathcal{O}_{y_u} = y_u H ^2 \bar{Q}_L \tilde{H} u_R + \text{h.c.}$	$\mathcal{O}_{y_d} = y_d H ^2 \bar{Q}_L H d_R + \text{h.c.}$	$\mathcal{O}_{y_e} = y_e H ^2 \bar{L}_L H e_R + \text{h.c.}$
$\mathcal{O}_R^u = (iH^\dagger \overleftrightarrow{D}_\mu H)(\bar{u}_R \gamma^\mu u_R)$	$\mathcal{O}_R^d = (iH^\dagger \overleftrightarrow{D}_\mu H)(\bar{d}_R \gamma^\mu d_R)$	$\mathcal{O}_R^e = (iH^\dagger \overleftrightarrow{D}_\mu H)(\bar{e}_R \gamma^\mu e_R)$
$\mathcal{O}_L^q = (iH^\dagger \overleftrightarrow{D}_\mu H)(\bar{Q}_L \gamma^\mu Q_L)$		
$\mathcal{O}_L^{(3)q} = (iH^\dagger \sigma^a \overleftrightarrow{D}_\mu H)(\bar{Q}_L \sigma^a \gamma^\mu Q_L)$		
$\mathcal{O}_{LL}^{(3)ql} = (\bar{Q}_L \sigma^a \gamma_\mu Q_L)(\bar{L}_L \sigma^a \gamma^\mu L_L)$		$\mathcal{O}_{LL}^{(3)l} = (\bar{L}_L \sigma^a \gamma_\mu L_L)(\bar{L}_L \sigma^a \gamma^\mu L_L)$

Table 1. Set of CP-even dimension-6 operators that defines our basis. We are including only the four-fermion operators that affect our observables. We omit dipole operators for fermions and $\mathcal{O}_R^{ud} = y_u^\dagger y_d (i\tilde{H}^\dagger \overleftrightarrow{D}_\mu H)(\bar{u}_R \gamma^\mu d_R)$ since they are suppressed by light fermion Yukawas under the MFV assumption. Also \mathcal{O}_{3G} is not included since it does not enter in our observables. The complete set of operators can be found in ref. [1].

of table 1 we assume a contraction of family indices i, j inside each parenthesis, e.g., $(\bar{Q}_L \gamma^\mu Q_L) = (\bar{Q}_L^i \gamma^\mu Q_L^i)$, and $y_u \bar{Q}_L u_R = y_u^{ij} \bar{Q}_L^i u_R^j$, as implied by the MFV assumption at the leading order in a Yukawa expansion. The top quark, having a large Yukawa coupling, could depart from the MFV assumption. For this reason we will also consider the impact of treating top operators separately.

CP-odd operators are not included in table 1, since they do not interfere with the SM contributions to the observables that we are considering. Their Wilson coefficients only enter quadratically in these processes and can then be neglected at the linear level that we are working. In section 4, however, we will briefly discuss their implications for TGC and $h \rightarrow V\bar{f}f$.

2.1 Experimental input values

We take the SM as defined by the 3 parameters g, g' and $v \simeq 246 \text{ GeV}$, that we relate with the well-measured values of the Fermi constant G_F as measured in muon decays, the fine-structure constant α_{em} , and the Z -boson mass m_Z . New physics, parametrized through the dimension-6 operators of table 1, affects these 3 input observables. The subset of relevant operators is

$$\Delta\mathcal{L}_{\text{input}} = \frac{c_T}{v^2} \mathcal{O}_T + \frac{c_V^\pm}{m_W^2} (\mathcal{O}_W + \mathcal{O}_B) + \frac{c_{LL}^{(3)l}}{v^2} \mathcal{O}_{LL}^{(3)l}, \quad (2.1)$$

where we have defined

$$c_V^\pm = \frac{1}{2}(c_W \pm c_B), \quad (2.2)$$

and $c_{W,B}$ are the Wilson coefficients of the operators $\mathcal{O}_{W,B}$. The orthogonal combination $\mathcal{O}_W - \mathcal{O}_B$ does not affect the input parameters. Notice that, instead of a common suppression scale Λ , we are suppressing the operators of eq. (2.1) by the weak scale, either m_W or v , meaning that we have absorbed the dependence on Λ into the Wilson coefficients. With this normalization the Wilson coefficients must be considered smaller than $O(1)$ for our expansion to make sense.

The modifications of the 3 input observables due to the dimension-6 operators are given by

$$\frac{\delta\alpha_{\text{em}}}{\alpha_{\text{em}}} = -2s_{\theta_W}^2 \hat{S}, \quad \frac{\delta m_Z^2}{m_Z^2} = -\hat{T} + 2s_{\theta_W}^2 \hat{S}, \quad \frac{\delta G_F}{G_F} = -2c_{LL}^{(3)l}, \quad (2.3)$$

where we have defined $s_{\theta_W} \equiv \sin\theta_W$ (and similarly for other trigonometric functions), being θ_W the weak mixing angle, and

$$\hat{S} = 2c_V^\dagger, \quad \hat{T} = c_T, \quad (2.4)$$

characterize the dominant contributions to the W/Z propagators [14, 15].

Our fits for electroweak observables are performed using a χ^2 analysis; the experimental data including correlations, are taken from refs. [16, 17], while the SM predictions, including the dominant loop corrections, are taken from refs. [18, 19] (we have checked that our results including only \hat{S} and \hat{T} agree with ref. [18] at the percent level). Henceforth when referring to SM predictions to observables we will always be considering predictions given as a function of these 3 input experimental values, G_F , α_{em} and m_Z .

2.2 Z -pole observables for leptons and the W mass

Among the most accurate experimental tests of the SM, stand the Z -pole observables measured at LEP-I/SLC, and the Tevatron measurement of the W mass. In a first set of constraints, we single out the leptonic observables that are measured at the per-mille level. The relevant dimension-6 operators that contribute to these quantities are, apart from $\Delta\mathcal{L}_{\text{input}}$,

$$\Delta\mathcal{L}_{\text{leptons}} = \frac{c_R^e}{v^2} \mathcal{O}_R^e. \quad (2.5)$$

It is important to notice that, being at the Z -pole, four-fermion operators can be neglected. LEP-I and SLC measurements afford only 3 observables for the lepton sector, that we can think of as $\Gamma_Z^{lL} \equiv \Gamma(Z \rightarrow \bar{l}_L l_L)$, $\Gamma_Z^{lR} \equiv \Gamma(Z \rightarrow \bar{l}_R l_R)$ and $\Gamma_Z^\nu \equiv \Gamma(Z \rightarrow \bar{\nu}\nu)$. These can be extracted from the correlated set of observables $A_l, R_l, \sigma_{\text{had}}^0, \Gamma_Z$ defined in eq. (A.6) of appendix A, where $\Gamma_Z^\nu = \Gamma_Z - \Gamma_Z^{\text{visible}}$ since we assume that there are no extra light degrees of freedom. The modifications of these quantities with respect to the SM predictions due to the dimension-6 operators are given by (see appendix A for details)

$$\frac{\delta\Gamma_Z^\nu}{\Gamma_Z^\nu} = \hat{T} - \frac{\delta G_F}{G_F}, \quad \frac{\delta\Gamma_Z^{lL}}{\Gamma_Z^{lL}} = \frac{1}{c_{2\theta_W}^2} \left(\hat{T} - \frac{\delta G_F}{G_F} - 4s_{\theta_W}^2 \hat{S} \right), \quad (2.6)$$

$$\frac{\delta\Gamma_Z^{lR}}{\Gamma_Z^{lR}} = -\frac{1}{c_{2\theta_W}^2} \left(\hat{T} - \frac{\delta G_F}{G_F} - 2\hat{S} \right) - \frac{c_R^e}{s_{\theta_W}^2}. \quad (2.7)$$

This gives 3 observables to constrain 4 Wilson coefficients. As an additional observable needed to constrain all 4 coefficients, we take the W mass, whose correction from the SM value is given by

$$\frac{\delta m_W}{m_W} = \frac{1}{2c_{2\theta_W}} \left[c_{\theta_W}^2 \hat{T} - s_{\theta_W}^2 \left(\frac{\delta G_F}{G_F} + 2\hat{S} \right) \right]. \quad (2.8)$$

Using this subset of observables, we can fit c_T , c_V^+ , c_R^e and $c_{LL}^{(3)l}$. Marginalizing over all but one coefficient at a time we find the 95% C.L. intervals:

$$\begin{aligned} c_T &\in [-5, 1] \times 10^{-3}, & c_V^+ &\in [-6, 0] \times 10^{-3}, \\ c_R^e &\in [-5, 0] \times 10^{-3}, & c_{LL}^{(3)l} &\in [-12, 2] \times 10^{-3}. \end{aligned} \quad (2.9)$$

Other four-lepton operators, apart from $\mathcal{O}_{LL}^{(3)l}$, can be constrained by LEP-II measurements of the cross-sections $e^+e^- \rightarrow l^+l^-$ [20]. These operators and their experimental constraints are completely orthogonal to our analysis and can be studied independently. We will not pursue them further.

2.3 Quarks

Experimental data for quark physics abounds. The most accurate measurements come, again, from LEP-I physics at the Z -pole. Assuming flavor-universality, the relevant dimension-6 operators for this type of physics are, apart from $\Delta\mathcal{L}_{\text{input}}$,

$$\Delta\mathcal{L}_{\text{quarks}} = \frac{c_L^q}{v^2} \mathcal{O}_L^q + \frac{c_L^{(3)q}}{v^2} \mathcal{O}_L^{(3)q} + \frac{c_R^u}{v^2} \mathcal{O}_R^u + \frac{c_R^d}{v^2} \mathcal{O}_R^d. \quad (2.10)$$

Eq. (2.10) contains 4 new parameters that need 4 new observables in order to be constrained. We take R_b , Γ_Z^{had} , A_b and A_c as defined in appendix A. They are affected by dimension-6 operators as

$$\frac{\delta\Gamma_Z^{\text{had}}}{\Gamma_Z^{\text{had}}} \simeq 0.7 c_L^q + 2.3 c_L^{(3)q} + 0.3 c_R^u - 0.2 c_R^d - 1.5 \delta s_{\theta_W}^2 + \hat{T} - \frac{\delta G_F}{G_F}, \quad (2.11)$$

$$\frac{\delta R_b}{R_b} \simeq 1.6 c_L^q - 0.05 c_L^{(3)q} - 0.3 c_R^u - 0.1 c_R^d + 0.2 \delta s_{\theta_W}^2, \quad (2.12)$$

$$\frac{\delta A_c}{A_c} \simeq -0.9 c_L^q + 0.9 c_L^{(3)q} - 2.3 c_R^u - 4.2 \delta s_{\theta_W}^2, \quad (2.13)$$

$$\frac{\delta A_b}{A_b} \simeq 0.1 c_L^q + 0.1 c_L^{(3)q} + 0.8 c_R^d - 0.6 \delta s_{\theta_W}^2, \quad (2.14)$$

where $\delta s_{\theta_W}^2$ is the correction to the effective weak mixing angle:

$$\delta s_{\theta_W}^2 = \frac{1}{c_{2\theta_W}} \left[s_{\theta_W}^2 \hat{S} - \frac{s_{2\theta_W}^2}{4} \left(\hat{T} - \frac{\delta G_F}{G_F} \right) \right]. \quad (2.15)$$

Only the first two observables, Γ_Z^{had} and R_b , are measured at the per-mille level [16] and can give strong constraints on $c_L^{(3)q}$ and c_L^q , since, accidentally, they enter respectively with large coefficients in eq. (2.11) and eq. (2.12).

The above hadronic observables can be complemented with information from low-energy measurements. One of the most relevant, due to its high accuracy, is the determination of the unitarity of the CKM matrix by KLOE and β -decays [21]. This result can be interpreted as a measurement of the Fermi constant in quark-lepton weak interactions, $G_F^{(q)}$, normalized to G_F as extracted from μ -decays:

$$\left(\frac{G_F^{(q)}}{G_F}\right)^2 = 0.9999(12) \quad \text{at 95 \% C.L.} \quad (2.16)$$

Dimension-six operators modify this SM prediction:

$$\delta \left(\frac{G_F^{(q)}}{G_F}\right)^2 = 4c_{LL}^{(3)l} + 2c_L^{(3)q} - 2c_{LL}^{(3)ql}, \quad (2.17)$$

which also includes a contribution from 4-fermion operators:

$$\Delta\mathcal{L}_{\text{quark-lepton}} = \frac{c_{LL}^{(3)ql}}{v^2} \mathcal{O}_{LL}^{(3)ql}. \quad (2.18)$$

Now, using recent LHC data [22] for the measurement of the high-energy differential cross-sections of $\bar{q}q \rightarrow \bar{l}\nu$, we can independently constrain $c_{LL}^{(3)ql}$, as shown in appendix B. This allows us to use eq. (2.17) to put an independent constraint on $c_L^{(3)q}$. Putting all the above information together and marginalizing (since the observables eqs. (2.11)–(2.14) are also sensitive to $\Delta\mathcal{L}_{\text{input}}$, we must include also the leptonic operators and observables and perform a global fit: the resulting constraints on the coefficients of $\Delta\mathcal{L}_{\text{input}}$ will be given in the conclusion), we obtain the 95% C.L. intervals:

$$\begin{aligned} c_L^q &\in [-1, 4] \times 10^{-3}, & c_L^{(3)q} &\in [-7, 4] \times 10^{-3}, \\ c_R^u &\in [-8, 0] \times 10^{-3}, & c_R^d &\in [-53, 1] \times 10^{-3}, \\ c_{LL}^{(3)ql} &\in [-2, 3] \times 10^{-3}. \end{aligned} \quad (2.19)$$

It is interesting to consider the case in which operators made of the top quark have different coefficients from those associated to operators made with the first two families. This is motivated by the largeness of the top Yukawa coupling that suggests that there can be large deviations from flavor-universality. We then add the following operator to our analysis:¹

$$\Delta\mathcal{L}_{\text{top}} = \frac{c_L^{+q3}}{v^2} \left(\mathcal{O}_L^{q3} + \mathcal{O}_L^{(3)q3} \right), \quad \text{where } c_L^{+q3} = \frac{1}{2}(c_L^{q3} + c_L^{(3)q3}), \quad (2.20)$$

that, due to the presence of b_L inside the 3rd family doublet Q_L , gives an extra contribution to the observables of eqs. (2.11)–(2.14):

$$\left. \frac{\delta\Gamma_Z^{\text{had}}}{\Gamma_Z^{\text{had}}} \right|_{q3} \simeq 0.5 c_L^{+q3}, \quad \left. \frac{\delta R_b}{R_b} \right|_{q3} \simeq 1.8 c_L^{+q3}, \quad \left. \frac{\delta A_b}{A_b} \right|_{q3} \simeq 0.1 c_L^{+q3}. \quad (2.21)$$

¹The other operators involving top-quarks, $(\mathcal{O}_L^{q3} - \mathcal{O}_L^{(3)q3})$ and \mathcal{O}_R^t , affect only top couplings to gauge bosons and their constraints arise from top physics at the LHC. We will not discuss them here.

We have now 5 coefficients in the quark sector, eq. (2.10) and eq. (2.20), but also 5 observables, those of eqs. (2.11)–(2.14) and eq. (2.16). KLOE data, which in the flavor-universal case was only used to slightly improve the numerical constraints, is now crucial to avoid unconstrained directions. The analysis gives:²

$$c_L^{+q3} \in [-7, 13] \times 10^{-3}. \quad (2.22)$$

Other four-quark operators, apart from eq. (2.17), can be constrained by LHC di-jet measurements [23, 24]. As for the lepton case, these bounds do not interfere with our analysis, neither are relevant for Higgs physics; for this reason they can be studied in a separate context.

2.4 Triple gauge-boson couplings

Extra important experimental data needed to further constrain our set of operators is the measure of the ZWW and γWW couplings. The best measurements still come from $e^+e^- \rightarrow W^+W^-$ at LEP-II, although LHC data is starting to be competitive. The dimension-6 operators involved in this process are, apart from eqs. (2.1) and (2.5),

$$\Delta\mathcal{L}_{\text{TGC}} = \frac{\kappa_{HV}^+}{m_W^2} (\mathcal{O}_{HW} + \mathcal{O}_{HB}) + \frac{c_V^- + \kappa_{HV}^-}{2m_W^2} \mathcal{O}_+ + \frac{\kappa_{3W}}{m_W^2} \mathcal{O}_{3W}, \quad (2.23)$$

where we have defined

$$\kappa_{HV}^\pm = \frac{1}{2}(\kappa_{HW} \pm \kappa_{HB}), \quad (2.24)$$

with $\kappa_{HW,HB}$ the Wilson coefficients of $\mathcal{O}_{HW,HB}$, and

$$\mathcal{O}_\pm \equiv (\mathcal{O}_W - \mathcal{O}_B) \pm (\mathcal{O}_{HW} - \mathcal{O}_{HB}). \quad (2.25)$$

Notice that in order to connect with the different experiments, we are, for convenience, working with the orthogonal combinations $\{\mathcal{O}_W + \mathcal{O}_B, \mathcal{O}_{HW} + \mathcal{O}_{HB}, \mathcal{O}_\pm\}$ instead of the original subset $\{\mathcal{O}_W, \mathcal{O}_B, \mathcal{O}_{HW}, \mathcal{O}_{HB}\}$ given in table 1. It is easy to show that the coefficients of \mathcal{O}_\pm are given by the combinations $(c_V^\pm \pm \kappa_{HV}^\pm)/2$. In eq. (2.23) we have not included the operator \mathcal{O}_- since it does not enter in TGC, and it is only relevant for Higgs physics as we will see later. Indeed, one finds the following contributions to the ZWW and γWW vertices [25]:

$$\begin{aligned} \Delta\mathcal{L}_{3V} &= ig\delta g_1^Z c_{\theta_W} Z^\mu (W^{-\nu} W_{\mu\nu}^+ - W^{+\nu} W_{\mu\nu}^-) + ig(\delta\kappa_Z c_{\theta_W} Z^{\mu\nu} + \delta\kappa_\gamma s_{\theta_W} A^{\mu\nu}) W_\mu^- W_\nu^+ \\ &\quad + \frac{ig}{m_W^2} (\lambda_Z c_{\theta_W} Z^{\mu\nu} + \lambda_\gamma s_{\theta_W} A^{\mu\nu}) W_\nu^{-\rho} W_{\rho\mu}^+, \end{aligned} \quad (2.26)$$

where $V_{\mu\nu} \equiv \partial_\mu V_\nu - \partial_\nu V_\mu$ for $V = W^\pm, Z, A$, and

$$\begin{aligned} \delta g_1^Z &= \frac{c_W + \kappa_{HW}}{c_{\theta_W}^2} = \frac{1}{c_{\theta_W}^2} (c_V^+ + c_V^- + \kappa_{HV}^- + \kappa_{HV}^+), \\ \delta\kappa_\gamma &= \kappa_{HW} + \kappa_{HB} = 2\kappa_{HV}^+, \quad \delta\kappa_Z = \delta g_1^Z - t_{\theta_W}^2 \delta\kappa_\gamma, \\ \lambda_Z &= \lambda_\gamma = \kappa_{3W}. \end{aligned} \quad (2.27)$$

²Introducing this extra parameter slightly alters the limits of eqs. (2.9) and (2.19), as we will show later.

Notice that out of the 5 quantities in eq. (2.26), only 3 are independent if one considers only dimension-6 operators; we take these to be δg_1^Z , $\delta\kappa_\gamma$ and λ_γ . This means that data from the differential cross-sections of $e^+e^- \rightarrow W^+W^-$ can only constrain 3 extra combinations of Wilson coefficients that corresponds to those of eq. (2.23): κ_{HV}^+ , $c_V^- + \kappa_{HV}^-$ and κ_{3W} .³ Unfortunately a three-parameter analysis has not been provided by the full LEP-II collaboration [17, 20], but only by DELPHI [26]. Nonetheless, a small value κ_{3W} is expected in most theories where the SM gauge bosons are assumed to be elementary above the BSM scale Λ . Neglecting this contribution (and then taking $\lambda_\gamma = 0$), we can use the two-parameter fit of g_1^Z and κ_γ from ref. [17], to obtain the 95% C.L. intervals:

$$c_V^- + \kappa_{HV}^- \in [-4.4, 6.6] \times 10^{-2}, \quad \kappa_{HV}^+ \in [-5.5, 3.9] \times 10^{-2}. \quad (2.28)$$

Including in the fit κ_{3W} does not change considerably our results; using DELPHI data [26] (which is however strongly correlated with ref. [17]) one finds results similar to eq. (2.28) and $\kappa_{3W} \in [-4, 7] \times 10^{-2}$.

3 Higgs physics

We now extend our analysis to Higgs physics. Apart from the operators already introduced before that could also affect Higgs physics, we have 8 CP-even dimension-six operators that give contributions only to Higgs physics and not to other SM processes. These are those operators that can be built from the Higgs modulus, i.e., $|H|^2$. In our basis, table 1, these operators are

$$\Delta\mathcal{L}_{\text{Higgs}} = \frac{c_H}{v^2}\mathcal{O}_H + \sum_{f=t,b,\tau} \frac{c_{y_f}}{v^2}\mathcal{O}_{y_f} + \frac{c_6}{v^2}\mathcal{O}_6 + \frac{\kappa_{BB}}{m_W^2}\mathcal{O}_{BB} + \frac{\kappa_{GG}}{m_W^2}\mathcal{O}_{GG} + \frac{c_V^- - \kappa_{HV}^-}{2m_W^2}\mathcal{O}_-, \quad (3.1)$$

which holds for one family, taken to be the 3rd one as it is the most relevant for Higgs physics. The presence of \mathcal{O}_- must be traced back [1] to the fact that the operator $\mathcal{O}_{WW} = g^2|H|^2W^{\mu\nu a}W_{\mu\nu}^a$, that can obviously affect only Higgs physics, when written in our basis is given by

$$\mathcal{O}_{WW} = 4\mathcal{O}_- + \mathcal{O}_{BB}. \quad (3.2)$$

Now, the operator \mathcal{O}_6 is in one-to-one correspondence with a deviation in the triple-Higgs coupling that has not yet been measured at the LHC. Therefore, we will not be considering it any longer. The 7 remaining operators of eq. (3.1) affect the main Higgs branching ratios, as well as Higgs production cross-sections at the LHC. In particular, we have that the 4 Wilson coefficients c_{y_b} , c_{y_τ} , κ_{BB} , κ_{HV}^- enter respectively in $BR(h \rightarrow bb)$, $BR(h \rightarrow \tau\tau)$, $BR(h \rightarrow \gamma\gamma)$, $BR(h \rightarrow Z\gamma)$, while κ_{GG} affects $\sigma(GG \rightarrow h)$; the coefficient c_H gives a contribution to the Higgs propagator and then enters universally in all Higgs processes (affecting, in particular, vector-boson fusion), while c_{y_t} modifies the $h\bar{t}t$ coupling that enters indirectly in $BR(h \rightarrow \gamma\gamma)$ and $BR(h \rightarrow Z\gamma)$ at the one-loop level [9], and also enters in

³Other Wilson coefficients entering in the process $e^+e^- \rightarrow W^+W^-$, such as those of eqs. (2.1) and (2.5), can be neglected since their constraints, derived in our previous analysis, are stronger than the ones we will be obtaining here.

the associated production process $pp \rightarrow h\bar{t}t$ that the LHC can become sensitive to in the near future. The Lagrangian for these processes is determined by

$$\mathcal{L}_h^{\text{SM}} + \Delta\mathcal{L}_h, \quad (3.3)$$

where

$$\mathcal{L}_h^{\text{SM}} = \frac{1}{2}(\partial_\mu h)^2 + (\sqrt{2}G_F)^{1/2} h \left[-m_f(\bar{f}_L f_R + \text{h.c.}) + m_Z^2 Z_\mu Z^\mu + 2m_W^2 W_\mu^+ W^{-\mu} \right] + \dots, \quad (3.4)$$

gives the SM contribution to single-Higgs processes written as a function of the physical masses, m_f , m_Z and $m_W^2 = m_Z^2 \left(1 + \sqrt{1 - 2^{3/2}\pi\alpha_{\text{em}}/m_Z^2 G_F}\right)/2$, while

$$\begin{aligned} \Delta\mathcal{L}_h = & \frac{c_H}{2}(\partial^\mu h)^2 + \frac{h}{v} \left[\delta g_{hff}(\bar{f}_L f_R + \text{h.c.}) + \delta g_{hZZ} Z_\mu Z^\mu + \delta g_{hWW} W_\mu^+ W^{-\mu} \right. \\ & \left. + 4\kappa_{BB} s_{\theta_W}^2 A_{\mu\nu} A^{\mu\nu} + 4\kappa_{GG} \frac{g_s^2}{g^2} G_{\mu\nu}^A G^{\mu\nu A} + 4t_{\theta_W} \kappa_{Z\gamma} Z_{\mu\nu} A^{\mu\nu} \right], \end{aligned} \quad (3.5)$$

gives the contributions proportional to the Wilson coefficients of $\Delta\mathcal{L}_{\text{input}}$, $\Delta\mathcal{L}_{\text{Higgs}}$, and also that of \mathcal{O}_+ in eq. (2.23) that enters in $hZ\gamma$. The first line of eq. (3.5) includes the corrections to the SM $h\bar{f}f$ and hVV interactions ($V = W, Z$) due to the Wilson coefficients entering in the input parameters, in the wave-function renormalization of V and in direct contributions arising from \mathcal{O}_T . They read

$$\frac{\delta g_{hZZ}}{g_{hZZ}^{\text{SM}}} = -c_T - \frac{\delta G_F}{2G_F}, \quad \frac{\delta g_{hWW}}{g_{hWW}^{\text{SM}}} = 2 \frac{\delta m_W}{m_W} - \frac{\delta G_F}{2G_F}, \quad \frac{\delta g_{hff}}{g_{hff}^{\text{SM}}} = -c_{y_f} - \frac{\delta G_F}{2G_F}. \quad (3.6)$$

With the exception of c_{y_f} , we know from the analysis of section 2 that these effects are small. The second line of eq. (3.5) gives the corrections to $h\gamma\gamma$, hGG and $hZ\gamma$ coupling, where

$$\kappa_{Z\gamma} = -\frac{1}{2}\kappa_{HV}^- - 2s_{\theta_W}^2 \kappa_{BB}. \quad (3.7)$$

The dominant modifications from eq. (3.5) to $BR(h \rightarrow \bar{f}f)$, $BR(h \rightarrow \gamma\gamma)$, $BR(h \rightarrow Z\gamma)$ and $\sigma(GG \rightarrow h)$ can be found in [9]. We confront these with a combination of ATLAS [27] and CMS [28–30] data (for technical details see ref. [31]). For this fit we also include the measurements of $pp \rightarrow h \rightarrow VV^*$ and vector-boson fusion $pp \rightarrow VVqq \rightarrow hqq$; there are extra contributions to these processes, as we will see in the next subsection (eqs. (3.9) and (3.10)), but these, we can advance, are found to be small and can be neglected here. Strong bounds can be found for Wilson coefficients entering at tree-level in the effective Higgs couplings to GG , $\gamma\gamma$ and $Z\gamma$, as they arise in the SM at the one-loop level. Indeed, marginalizing over c_H and c_{y_f} , we obtain at the 95% C.L.

$$\kappa_{GG} \in [-0.8, 0.8] \times 10^{-3}, \quad \kappa_{BB} \in [-1.3, 1.8] \times 10^{-3}, \quad \kappa_{Z\gamma} \in [-6, 12] \times 10^{-3}. \quad (3.8)$$

Notice that the constraints on $g_S^2 \kappa_{GG}$ are comparable to those on $g'^2 \kappa_{BB}$, and that the bound on the $hZ\gamma$ coupling is quite strong even though the experimental data is only sensitive to $BR(h \rightarrow Z\gamma)$ that are ~ 10 above the SM value [32, 33]. For the Wilson

coefficients c_H and c_{y_f} we do not find significant bounds after marginalizing over the others (for this reason we do not report them). Although this implies that there is room for BSM here, we recall that these coefficients cannot be larger than one if we want our expansion (in higher-dimensional operators) to be reliable.

This ends our analysis for getting the most important constraints on the operators of table 1. Our results can be particularly useful to bound possible new physics effects in future new Higgs measurements. We show this below with the example of the $h \rightarrow V \bar{f} f$ decay.

3.1 New physics effects in $h \rightarrow V \bar{f} f$

The decays $h \rightarrow V \bar{f} f$ ($V = W, Z$) are potentially much richer than two-body decays, since the different differential partial-widths can give in principle extra information on BSM contributions [7, 34–39]. Nevertheless, as we will show, most of the new information that we could extract from measuring the various differential partial-widths of the decay $h \rightarrow V \bar{f} f$ is already constrained by other experiments.

Contributions to $h \rightarrow V \bar{f} f$ can come from corrections to hVV vertices and contact-interactions $hV\bar{f}f$. Apart from the contributions given in eq. (3.3), we have

$$\begin{aligned} \Delta\mathcal{L}_{hVV} = & 2\frac{h}{v} \left[\hat{c}_W (W_\mu^- \mathcal{D}^{\mu\nu} W_\nu^+ + \text{h.c.}) + \hat{c}_Z Z_\mu \mathcal{D}^{\mu\nu} Z_\nu + (\hat{c}_W - \hat{c}_B) t_{\theta_W} Z_\mu \mathcal{D}^{\mu\nu} A_\nu \right] \\ & - 2\frac{h}{v} \left[c_{WW} W^{+\mu\nu} W_{\mu\nu}^- + c_{ZZ} Z^{\mu\nu} Z_{\mu\nu} \right], \end{aligned} \quad (3.9)$$

$$\Delta\mathcal{L}_{hVff} = \frac{h}{v} \sum_{f=f_L, f_R} \left[g_{hWff'} W_\mu \bar{f} \gamma^\mu f' + g_{hZff} Z_\mu \bar{f} \gamma^\mu f \right], \quad (3.10)$$

where $\mathcal{D}_{\mu\nu} = \partial_\mu \partial_\nu - \square \eta_{\mu\nu}$ and

$$\hat{c}_W = c_W + \kappa_{HW}, \quad \hat{c}_Z = \hat{c}_W + \hat{c}_B t_{\theta_W}^2, \quad \hat{c}_B = c_B + \kappa_{HB}, \quad (3.11)$$

$$c_{WW} = \kappa_{HW}, \quad c_{ZZ} = \frac{1}{2}(\kappa_{HW} + \kappa_{HB} t_{\theta_W}^2) - 2\frac{s_{\theta_W}^4}{c_{\theta_W}^2} \kappa_{BB}. \quad (3.12)$$

Eq. (3.10) gives the contributions to the contact $hV\bar{f}f$ vertices that is found to be correlated with those to the $V\bar{f}f$ vertices:

$$g_{hZff} = \frac{2}{v} \delta g_Z^f \quad \text{and} \quad g_{hWff'} = \frac{2}{v} \delta g_W^f, \quad (3.13)$$

where δg_Z^f and δg_W^f are given respectively in eqs. (A.3) and (A.11) of appendix A.

The CP-even part of the total amplitude for the process $h \rightarrow V \bar{f} f$ can be written as⁴

$$\mathcal{M}(h \rightarrow V J_f) = (\sqrt{2}G_F)^{1/2} \epsilon^{*\mu}(q) J_f^V{}^\nu(p) \left[A_f^V \eta_{\mu\nu} + B_f^V (p \cdot q \eta_{\mu\nu} - p_\mu q_\nu) \right], \quad (3.14)$$

where q and p are respectively the total 4-momentum of V and the fermion pair in the J_f^V current ($J_{f_{L,R}}^\mu = \bar{f}_{L,R} \gamma^\mu f_{L,R}$), ϵ^μ is the polarization 4-vector of V , and we have defined

$$A_f^V = a_f^V + \hat{a}_f^V \frac{p^2 + m_V^2}{p^2 - m_V^2}, \quad B_f^V = b_f^V \frac{1}{p^2 - m_V^2} + \hat{b}_f^V \frac{1}{p^2} \quad (\hat{b}_f^V = 0 \text{ for } V = W). \quad (3.15)$$

⁴We neglect terms proportional to the light fermion masses (see however ref. [38]). Also we omit a term proportional to $C_f^W \epsilon_{\mu\nu\alpha\beta} p^\alpha q^\beta$ that could be CP-even if C_f^W is pure imaginary. None of the Wilson coefficients of the dimension-6 operators contribute to this term at tree-level.

The above coefficients are in one-to-one correspondence with the coefficients of the Lagrangians eqs. (3.5), (3.9) and (3.10):

$$\begin{aligned}
 a_f^Z &= -g_Z^f \left(1 + \frac{\delta g_{hZZ}}{m_Z^2} \right) + 2eQ_f(\hat{c}_W - \hat{c}_B)t_{\theta_W} + v g_{hZff}, & a_f^W &= -g_W^f \left(1 + \frac{\delta g_{hWW}}{2m_W^2} \right) + v g_{hWff'}, \\
 \hat{a}_f^Z &= g_Z^f \left(1 + \frac{\delta g_{hZZ}}{m_Z^2} + 2\hat{c}_Z \right), & \hat{a}_f^W &= g_W^f \left(1 + \frac{\delta g_{hWW}}{2m_W^2} + 2\hat{c}_W \right), \\
 b_f^Z &= 8g_Z^f c_{ZZ}, & b_f^W &= 4g_W^f c_{WW}, \\
 \hat{b}_f^Z &= -8eQ_f t_{\theta_W} \kappa_{Z\gamma}, & &
 \end{aligned} \tag{3.16}$$

where we have not included the universal contribution from c_H that drops when calculating BR . All contributions to the 7 quantities in eq. (3.16) correspond to Wilson coefficients that have already been constrained by other experiments. Indeed, the terms proportional to δg_{hVV} of eq. (3.6) and the universal part of the shifts in g_V^f , given in eqs. (A.4) and (A.11), are constrained, as we have shown, at the per-mil level, and can be readily neglected. Similarly, κ_{BB} that appears in eq. (3.12) is constrained from $BR(h \rightarrow \gamma\gamma)$. The contributions proportional to g_{hZff} that depends, in particular, on $c_R^{u,d}$, are also constrained but, as shown in eq. (2.19), only at the per-cent level. Nevertheless, their effects in total amplitudes squared, when summed over the different flavors, can be constrained much more. This is due to the interesting relation

$$\sum_{f=\text{quarks}} g_Z^f g_{hZff} = \frac{2}{v} \sum_{f=\text{quarks}} g_Z^f \delta g_Z^f = \frac{1}{v} \frac{\delta \Gamma_Z^{\text{had}}}{\Gamma_Z^{\text{had}}} \sum_{f=\text{quarks}} (g_Z^f)^2 + O(\hat{S}, \hat{T}, \delta G_F/G_F), \tag{3.17}$$

where in the first equality we have used eq. (3.13) and in the second we have used the hadronic part of eq. (A.9) that is true up to pieces proportional to \hat{S}, \hat{T} and $\delta G_F/G_F$. Since all quantities on the right-hand side of eq. (3.17) are very well measured (per-mille level), we obtain a constraint on the effect of g_{hZff} at this level. Neglecting these terms, we are left with only 3 Wilson coefficients, c_V^- and κ_{HV}^\pm , that are constrained by the measurements of the $Z(\gamma)WW$ and $hZ\gamma$ couplings and whose bounds are not so strong. Using eq. (2.27) and eq. (3.7), we can relate the 7 coefficients of eq. (3.16) with the 3 experimental values of δg_1^Z , $\delta \kappa_\gamma$ and $\kappa_{Z\gamma}$. The result is

$$\begin{aligned}
 a_f^Z &\simeq -g_Z^{f,\text{SM}} + 2eQ_f (s_{2\theta_W} \delta g_1^Z - t_{\theta_W} \delta \kappa_\gamma), & a_f^W &\simeq -g_W^{f,\text{SM}}, \\
 \hat{a}_f^Z &\simeq g_Z^{f,\text{SM}} (1 + 2c_{2\theta_W} \delta g_1^Z + 2t_{\theta_W}^2 \delta \kappa_\gamma), & \hat{a}_f^W &= g_W^{f,\text{SM}} (1 + 2c_{\theta_W}^2 \delta g_1^Z), \\
 b_f^Z &\simeq \frac{2g_Z^{f,\text{SM}}}{c_{\theta_W}^2} (\delta \kappa_\gamma - 4c_{2\theta_W} \kappa_{Z\gamma}), & b_f^W &\simeq 2g_W^{f,\text{SM}} (\delta \kappa_\gamma - 4\kappa_{Z\gamma}), \\
 \hat{b}_f^Z &\simeq -8eQ_f t_{\theta_W} \kappa_{Z\gamma}. & &
 \end{aligned} \tag{3.18}$$

We can use the experimental constraints on δg_1^Z , $\delta \kappa_\gamma$ and $\kappa_{Z\gamma}$ to put a constraint on the

size of these 7 quantities. For example, for the case of $h \rightarrow Z\bar{l}l$, we find

$$\begin{aligned} \frac{\delta a_{l_L}^Z}{a_{l_L}^Z} &\in [-0.2, 0.1], & \frac{\delta \hat{a}_{l_L}^Z}{\hat{a}_{l_L}^Z} &\in [-8, 7] \times 10^{-2}, & b_{l_L}^Z &\in [-2, 5] \times 10^{-2}, & \hat{b}_{l_L}^Z &\in [-2, 5] \times 10^{-2}, \\ \frac{\delta a_{l_R}^Z}{a_{l_R}^Z} &\in [-0.2, 0.3], & \frac{\delta \hat{a}_{l_R}^Z}{\hat{a}_{l_R}^Z} &\in [-8, 7] \times 10^{-2}, & b_{l_R}^Z &\in [-3, 2] \times 10^{-2}, & \hat{b}_{l_R}^Z &\in [-2, 5] \times 10^{-2}. \end{aligned}$$

Although the allowed range in $a_{l_{L,R}}^Z$ is quite large, we notice that their impact on the total amplitude, when summed over lepton chiralities, is much smaller, $2 \sum_{l=L,R} g_Z^l a_l^Z / \sum_{l=L,R} (g_Z^l)^2 \in [-6, 4] \times 10^{-2}$.

It is interesting to notice that in the limit $g' \rightarrow 0$ the result of eq. (3.18) is custodial invariant, i.e., one finds equal corrections for Z and W . This is because, for $g' \rightarrow 0$, the only Wilson coefficients breaking the custodial symmetry are c_T and $c_{L,R}^f$ [1] that, being constrained at the per-mille, have been dropped from eq. (3.18). We then find that the test of the custodial symmetry used at LHC [40] defined as

$$\lambda_{WZ}^2 \equiv \frac{\Gamma(h \rightarrow WW^{(*)})}{\Gamma^{\text{SM}}(h \rightarrow WW^{(*)})} \frac{\Gamma^{\text{SM}}(h \rightarrow ZZ^{(*)})}{\Gamma(h \rightarrow ZZ^{(*)})}, \quad (3.19)$$

is constrained by

$$\begin{aligned} \lambda_{WZ}^2 - 1 &\simeq s_{\theta_W}^2 [0.9c_W - 2.6c_B + 3\kappa_{HW} - 3.9\kappa_{HB}] \\ &\simeq 0.6\delta g_1^Z - 0.5\delta\kappa_\gamma - 1.6\kappa_{Z\gamma} \in [-6, 8] \times 10^{-2}, \end{aligned} \quad (3.20)$$

where the numerical values of the first line have been taken from [5]. We see that eq. (3.20) puts a bound on λ_{WZ} much stronger than the present direct experimental limit given by [27]:⁵ $(\lambda_{WZ} - 1) \in [-0.5, 0.1]$.

Along the lines presented here we could also study the corrections from Wilson coefficients to the Higgs production modes $\bar{f}f \rightarrow Vh$ and $VV \rightarrow h$ that we could similarly show that are constrained by our previous analysis.

4 CP-odd operators

For completeness, we show here how CP-odd operators enter in TGC and in the process $h \rightarrow V\bar{f}f$, and how they can be related. These operators are⁶

$$\begin{aligned} \mathcal{O}_{B\tilde{B}} &= g'^2 |H|^2 B_{\mu\nu} \tilde{B}^{\mu\nu}, & \mathcal{O}_{G\tilde{G}} &= g_s^2 |H|^2 G_{\mu\nu}^A \tilde{G}^{A\mu\nu}, \\ \mathcal{O}_{H\tilde{W}} &= ig (D^\mu H)^\dagger \sigma^a (D^\nu H) \tilde{W}_{\mu\nu}^a, & \mathcal{O}_{H\tilde{B}} &= ig' (D^\mu H)^\dagger (D^\nu H) \tilde{B}_{\mu\nu}, \\ \mathcal{O}_{3\tilde{W}} &= \frac{1}{3!} g \epsilon_{abc} W_\mu^{a\nu} W_{\nu\rho}^b \tilde{W}^{c\rho\mu}, \end{aligned} \quad (4.1)$$

⁵The experimental bound on λ_{WZ} is extracted not only using eq. (3.19) but also considering custodial breaking effects in vector-boson fusion. The impact of these latter effects is however negligible.

⁶A CP-odd operator involving 3 gluon field-strengths and the operators $iy_f |H|^2 \bar{f}_L H f_R + \text{h.c.}$ complete the list of CP-odd operators; since they do not enter in the observables discussed here, they have been omitted.

where $\tilde{V}^{\mu\nu} = \epsilon^{\mu\nu\rho\sigma} V_{\rho\sigma}/2$ for $V = W, B, G$. These operators affect TGC as [25]:

$$\Delta\mathcal{L}_{3\tilde{V}} = ig\left(\delta\tilde{\kappa}_Z c_{\theta_W} \tilde{Z}^{\mu\nu} + \delta\tilde{\kappa}_\gamma s_{\theta_W} \tilde{A}^{\mu\nu}\right) W_\mu^- W_\nu^+ + \frac{ig}{m_W^2}\left(\tilde{\lambda}_Z c_{\theta_W} \tilde{Z}^{\mu\nu} + \tilde{\lambda}_\gamma s_{\theta_W} \tilde{A}^{\mu\nu}\right) W_\nu^- W_{\rho\mu}^+,$$

where

$$\delta\tilde{\kappa}_\gamma = \kappa_{H\tilde{W}} + \kappa_{H\tilde{B}}, \quad \delta\tilde{\kappa}_Z = -t_{\theta_W}^2 \delta\tilde{\kappa}_\gamma, \quad \tilde{\lambda}_Z = \tilde{\lambda}_\gamma = \kappa_{3\tilde{W}}. \quad (4.2)$$

Although the experimental collaborations have not devoted explicit searches for these types of structures, constraints on the above coefficients can be obtained in a near future. For example, ref. [41] finds that with 10 fb^{-1} the 14 TeV LHC can potentially obtain $|\delta\tilde{\kappa}_\gamma| \lesssim 0.06$.

The contribution from CP-odd operators to Higgs physics can be read from the Lagrangian

$$\begin{aligned} \Delta\mathcal{L}_{hV\tilde{V}} = & -2\frac{h}{v} \left[c_{W\tilde{W}} W^{+\mu\nu} \tilde{W}_{\mu\nu}^- + c_{Z\tilde{Z}} Z^{\mu\nu} \tilde{Z}_{\mu\nu} - 2t_{\theta_W} \tilde{\kappa}_{Z\gamma} A^{\mu\nu} \tilde{Z}_{\mu\nu} \right. \\ & \left. - 2\kappa_{B\tilde{B}} s_{\theta_W}^2 A_{\mu\nu} \tilde{A}^{\mu\nu} - 2\kappa_{G\tilde{G}} \frac{g_s^2}{g^2} G_{\mu\nu}^A \tilde{G}^{\mu\nu A} \right], \end{aligned} \quad (4.3)$$

where, quite analogously to the CP-even case, we have

$$c_{W\tilde{W}} = \kappa_{H\tilde{W}}, \quad c_{Z\tilde{Z}} = \frac{1}{2}(\kappa_{H\tilde{W}} + \kappa_{H\tilde{B}} t_{\theta_W}^2) - 2\frac{s_{\theta_W}^4}{c_{\theta_W}^2} \kappa_{B\tilde{B}}, \quad \tilde{\kappa}_{Z\gamma} = \frac{1}{4}(\kappa_{H\tilde{B}} - \kappa_{H\tilde{W}}) - 2s_{\theta_W}^2 \kappa_{B\tilde{B}}. \quad (4.4)$$

Finally, the CP-odd contribution to the $h \rightarrow V\bar{f}f$ amplitude is given by

$$\mathcal{M}(h \rightarrow V J_f) = (\sqrt{2}G_F)^{1/2} \epsilon^{*\mu}(q) J_f^{V\nu}(p) \left[C_f^V \epsilon_{\mu\nu\alpha\beta} p^\alpha q^\beta \right], \quad (4.5)$$

where we find [42]

$$C_f^V = c_f^V \frac{1}{p^2 - m_V^2} + \hat{c}_f^V \frac{1}{p^2} \quad (\hat{c}_f^V = 0 \text{ for } V = W), \quad (4.6)$$

with

$$c_f^Z = 8g_Z^f c_{Z\tilde{Z}}, \quad c_f^W = 4g_W^f c_{W\tilde{W}}, \quad (4.7)$$

$$\hat{c}_f^Z = -8eQ_f t_{\theta_W} \tilde{\kappa}_{Z\gamma}. \quad (4.8)$$

Using eq. (4.2) and eq. (4.4), we find the relations

$$c_f^Z = \frac{2g_Z^f}{c_{\theta_W}^2} (\delta\tilde{\kappa}_\gamma - 4c_{2\theta_W} \tilde{\kappa}_{Z\gamma} - 2s_{2\theta_W}^2 \kappa_{B\tilde{B}}), \quad c_f^W = 2g_W^f (\delta\tilde{\kappa}_\gamma - 4\tilde{\kappa}_{Z\gamma} - 8s_{\theta_W}^2 \kappa_{B\tilde{B}}). \quad (4.9)$$

To our knowledge, the only CP-odd observables measured at present are those in the decay $h \rightarrow ZZ^* \rightarrow 4l$ [29], but only very weak constraints can be extracted for the respective Wilson coefficients. Nevertheless, the experimental bound on $BR(h \rightarrow Z\gamma)$ and $BR(h \rightarrow \gamma\gamma)$ can also be used to constrain the squared of $\tilde{\kappa}_{Z\gamma}$ and $\kappa_{B\tilde{B}}$.⁷ One obtains similar constraints to their CP-even counterparts, since the interference term involves the SM contribution that is one-loop suppressed. Neglecting $\tilde{\kappa}_{Z\gamma}$ and $\kappa_{B\tilde{B}}$ in eq. (4.9), one obtains a one-to-one relation between the CP-odd contributions to $h \rightarrow V\bar{f}f$ and the TGC parameter $\delta\tilde{\kappa}_\gamma$ that can be measured in the near future.

⁷The CP-odd and CP-even contributions add quadratically in the Higgs branching ratios, giving then independent bounds on the corresponding Wilson coefficients.

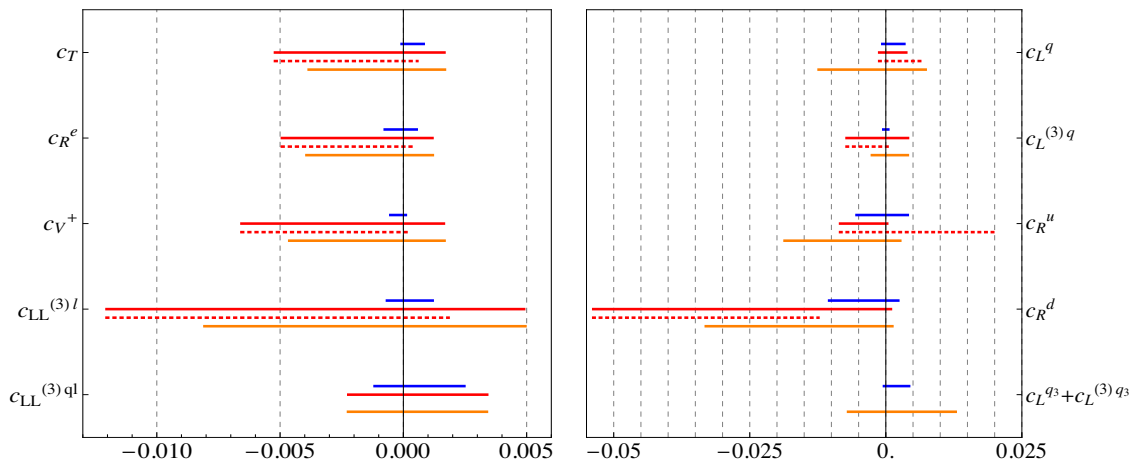


Figure 1. Bounds on the Wilson coefficients of eqs. (2.1), (2.5), (2.10) and (2.18), from LEP-I, SLC and KLOE data (and the necessary information from $pp \rightarrow l\bar{\nu}$ at the LHC — see appendix B). Solid red lines denote the 95% C.L. intervals after marginalizing over all other parameters (dashed lines do not include KLOE data); blue lines are obtained by setting all other coefficients to zero. The global fit including the operator of eq. (2.20) corresponds to the orange lines.

5 Conclusions

We have made a first step towards a complete SM fit, focusing here on EWSB effects in gauge-bosons and Higgs physics. We have characterized all possible deformations from SM physics using the Wilson coefficients of the independent dimension-6 operators, the relevant ones for our analysis given in table 1. Assuming flavor-universality (but also including the top-quark operators separately) and taking as the input data m_Z , G_F , α_{em} and m_f , m_h , we have grouped the operators according to their impact on the different experimental data.

In a first group we have the operators that can affect the gauge-boson propagators and their couplings to fermions. These receive strong constraints from well-measured quantities, mainly the Z -pole observables at LEP-I and SLC, the W mass at Tevatron, together with the check of the CKM unitarity from low-energy data. The constraints on the corresponding Wilson coefficients from a global fit are summarized in figure 1. To stress the correlation between the various operators, we compare the 95% C.L. contours obtained by marginalizing all other coefficients (in red) with the contours obtained by setting all other coefficients to zero (in blue). More data, such as the low-energy determination of the ν -nucleon and ν - e scattering or atomic parity violation experiments, could be added to our analysis. Nevertheless, due to their poorer resolution, we do not expect these data to affect substantially our results.

In a second group we have operators affecting triple gauge-boson vertices, that contrary to the previous group, receive milder constraints. Our basis is suitable for treating separately these two groups, while other bases, such as the one of ref. [11] used in the fit of ref. [7], makes this separation more difficult, due to strong correlations between bounds on different Wilson coefficients. We have 3 (combinations of) Wilson coefficients parametrizing these deviations, given in eq. (2.23). Using the 2-parameter fit of LEP-II [17] where

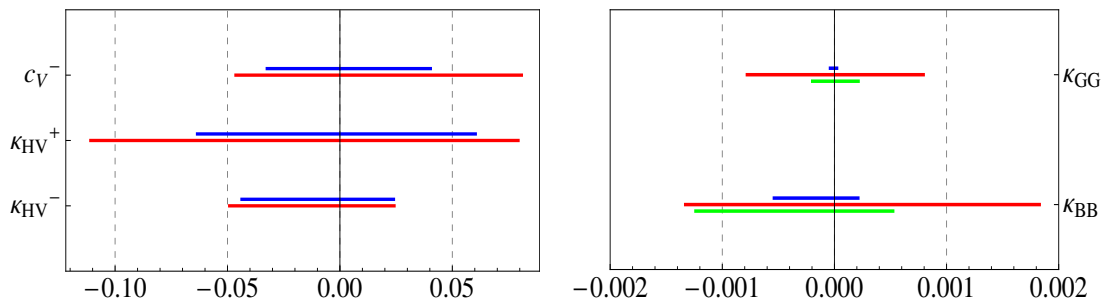


Figure 2. Left: bounds on some of the Wilson coefficients of eq. (2.23) from LEP-II TGC data and $BR(h \rightarrow Z\gamma)$. Right: bounds on some of the Wilson coefficients of eq. (3.1) from measurements of the Higgs BR at the LHC. Color code as in figure 1, with green lines including the theoretical prior on the coefficients c_H and c_{y_f} that cannot introduce modifications larger than 50% on SM predictions.

λ_γ (and therefore κ_{3W}) is neglected, we have presented bounds on two of them, eq. (2.28). Even if we allow for a nonzero λ_γ , LEP-II data is expected to be able to constrain all the 3 coefficients at the per-cent level [26], but a combined three-parameter fit is still not available. In figure 2 we show our results as bounds on the two coefficients c_V^- and κ_{HV}^+ (once the data on $BR(h \rightarrow Z\gamma)$ is also used to constrain the coefficient κ_{HV}^-).

Finally, there are the operators that only affect Higgs physics. These are 8 CP-even operators (for a single family) that can affect, independently, the Higgs BR to fermions, $\gamma\gamma$ and $Z\gamma$, as well as production cross-sections both in the gluon-gluon and vector-boson fusion modes (and also the triple-Higgs coupling). In figure 2 we present the model-independent constraints on the Wilson coefficients entering in the $h\gamma\gamma$ and hGG effective-couplings. These are very severe bounds as the BSM contributions enter at the tree-level and can then compete with the SM contributions arising at the one-loop level. Constraints on other Wilson coefficients, c_H and c_{y_f} , are at present very mild and therefore not shown. The Wilson coefficient κ_{HV}^- enters in $BR(h \rightarrow Z\gamma)$ and is constrained, as we said, only from the present experimental bound $BR(h \rightarrow Z\gamma) \lesssim 10 \times BR(h \rightarrow Z\gamma)_{SM}$ [32, 33], leaving still room for large deviations from the SM.

Having determined all Wilson coefficients, any other (Higgs) process will depend on the above physics and therefore their BSM effects can be indirectly constrained. In particular, we have studied $h \rightarrow V\bar{f}f$. Neglecting corrections tightly bounded from the first group of observables, we have shown that the different form-factors of the $h \rightarrow V\bar{f}f$ amplitude are related with TGC and $BR(h \rightarrow Z\gamma)$. These latter are at present already constrained, and in the future LHC can considerably improve these bounds, being more competitive than looking for deviations in the $h \rightarrow V\bar{f}f$ decays. Similar arguments apply to CP-odd operators.

Possible deviations in the approximate relations derived here, eq. (3.18), eq. (3.20) and eq. (4.9), would imply the breakdown of our assumptions, hinting possibly towards a non-linear realization of EWSB, in which h is not part of a Higgs doublet and no connection between Higgs and gauge-boson physics can be made.

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A Corrections to $V\bar{f}f$ couplings and Z -pole observables

The couplings between the Z -boson and fermions are altered as

$$g_Z^f = g_Z^{f,\text{SM}} + \delta g_Z^f + \delta g_Z^{f,\text{uni}}, \quad (\text{A.1})$$

where

$$g_Z^{f,\text{SM}} = m_Z \left(\sqrt{2} G_F \right)^{1/2} \left(T_{L,f}^3 - Q_f s_{\theta_W}^2 \right), \quad (\text{A.2})$$

is the SM value, $T_{L,f}^3$ and Q_f are respectively the weak-isospin and charge of the fermion $f = \{f_L, f_R\}$, and $s_{\theta_W}^2 \equiv \left(1 - \sqrt{1 - 2^{3/2} \pi \alpha_{\text{em}} / m_Z^2 G_F} \right) / 2$. We have divided the corrections into a fermion-specific part δg_Z^f ,

$$\delta g_Z^{f_L} = m_Z \left(\sqrt{2} G_F \right)^{1/2} \left(T_{L,f}^3 c_L^{(3)f} - \frac{c_L^f}{2} \right), \quad \delta g_Z^{f_R} = -m_Z \left(\sqrt{2} G_F \right)^{1/2} \frac{c_R^f}{2}, \quad (\text{A.3})$$

(notice that for leptons $\delta g_Z^{l_L} = 0$ in our basis) and a universal part,

$$\delta g_Z^{f,\text{uni}} = -m_Z \left(\sqrt{2} G_F \right)^{1/2} Q_f \delta s_{\theta_W}^2 + \frac{g_Z^f}{2} \left(\hat{T} - \frac{\delta G_F}{G_F} \right), \quad (\text{A.4})$$

which includes modifications to the wave-function of the Z -boson, $\delta \Pi'_Z(m_Z^2) = 2 \hat{S} s_{\theta_W}^2$, corrections to the input parameters, and contributions to $\Pi_{\gamma Z}(m_Z^2)$ through

$$\delta s_{\theta_W}^2 = \frac{1}{c_{2\theta_W}} \left[s_{\theta_W}^2 \hat{S} - \frac{s_{2\theta_W}^2}{4} \left(\hat{T} - \frac{\delta G_F}{G_F} \right) \right]. \quad (\text{A.5})$$

This is enough to compute the corrections to the observables that we use in the fit (the uncorrelated subsets $\{A_l, R_l, \sigma_{\text{had}}^0, \Gamma_Z\}$ and $\{R_b, R_c, A_{FB}^{0,b}, A_{FB}^{0,c}, A_b, A_c\}$ from ref. [16]):

$$A_l \equiv \frac{(g_Z^{l_L})^2 - (g_Z^{l_R})^2}{(g_Z^{l_L})^2 + (g_Z^{l_R})^2}, \quad R_l \equiv \frac{\Gamma_Z^{\text{had}}}{\Gamma_Z^l}, \quad \sigma_{\text{had}}^0 \equiv \frac{12\pi}{m_Z^2} \frac{\Gamma_Z^l \Gamma_Z^{\text{had}}}{\Gamma_Z^2}, \quad (\text{A.6})$$

and

$$R_q = \frac{(g_Z^{qL})^2 + (g_Z^{qR})^2}{3[(g_Z^{dL})^2 + (g_Z^{dR})^2] + 2[(g_Z^{uL})^2 + (g_Z^{uR})^2]}, \quad A_q = \frac{(g_Z^{qL})^2 - (g_Z^{qR})^2}{(g_Z^{qL})^2 + (g_Z^{qR})^2}, \quad A_{FB}^{0,q} = \frac{3}{4} A_q A_l, \quad (\text{A.7})$$

where $q = b, c$ at LEP. The partial widths are defined as,

$$\begin{aligned} \Gamma_Z^l &= \frac{m_Z}{24\pi} [(g_Z^{lL})^2 + (g_Z^{lR})^2], & \Gamma_Z^{\text{had}} &= \frac{m_Z}{24\pi} \left(3N_c [(g_Z^{dL})^2 + (g_Z^{dR})^2] + 2N_c [(g_Z^{uL})^2 + (g_Z^{uR})^2] \right), \\ \Gamma_Z^\nu &= \frac{m_Z}{24\pi} (g_Z^\nu)^2, & \Gamma_Z &= \Gamma_Z^{\text{had}} + 3\Gamma_Z^l + 3\Gamma_Z^\nu. \end{aligned} \quad (\text{A.8})$$

Notice that the universal corrections due to $\delta\Pi'_Z(m_Z^2)$ cancel out from the observables eq. (A.6), but enter in the total width (and similarly for the partial widths) which is corrected as

$$\frac{\delta\Gamma_Z}{\Gamma_Z} = 2 \frac{\sum_f g_Z^f (\delta g_Z^{f,\text{uni}} + \delta g_Z^f)}{\sum_f (g_Z^f)^2}. \quad (\text{A.9})$$

Finally, the couplings of W -bosons to fermions are modified as

$$g_W^f = g_W^{f,\text{SM}} + \delta g_W^f + \delta g_W^{f,\text{uni}}, \quad (\text{A.10})$$

where

$$g_W^{f,\text{SM}} = \frac{m_W G_F^{1/2}}{2^{1/4}}, \quad \frac{\delta g_W^f}{g_W^f} = c_L^{(3)f}, \quad \frac{\delta g_W^{f,\text{uni}}}{g_W^f} = \frac{\delta m_W}{m_W} - \frac{\delta G_F}{2G_F}. \quad (\text{A.11})$$

B $q\bar{q}'l\bar{\nu}$ contact-interactions at the LHC

Since the fermion-fermion scattering amplitude mediated by contact-interactions grows with energy, these can be tested with accuracy at the LHC. We are interested here to put a bound on the 4-fermion operator $\mathcal{O}_{LL}^{(3)ql}$. This operator affects in particular the cross-section of $pp \rightarrow l\bar{\nu}$ that has been measured at the LHC. Nevertheless, the only LHC analysis [22] has been on the helicity non-conserving (HNC) operator $(Q_L^r \bar{u}_R) \epsilon_{rs} (L_L^s \bar{e}_R)$, which also modifies $pp \rightarrow l\bar{\nu}$. Contrary to $\mathcal{O}_{LL}^{(3)ql}$, however, the HNC operator does not interfere with the SM contribution, so that the results of ref. [22] cannot be used for $\mathcal{O}_{LL}^{(3)ql}$ and a dedicated analysis is necessary.

A signal sample needs to be considered for every value of $c_{LL}^{(3)ql}$ and then compared with the data. Other operators entering in $pp \rightarrow l\bar{\nu}$, such as the HNC operator, are suppressed by small Yukawa couplings under the MFV assumption, that we consider here, and can be neglected. We simulate the effects of $\mathcal{O}_{LL}^{(3)ql}$ by integrating out a heavy W' -boson, implemented with FeynRules [43] and generate parton-level events with MadGraph [44]. We divide the signal and the data in 3 regions, according to the transverse mass of electron and neutrino: $M_T \in [1, 1.5]$ TeV, $M_T \in [1.5, 2]$ TeV and $M_T > 2$ TeV, which we treat as statistically independent. Ref. [22] reports, respectively, the observation of 22, 0 and 1 events in these regions. We compute C.L. contours, assuming a Poissonian distribution around the signal+background and treating errors as nuisances (an estimated 5% systematic error

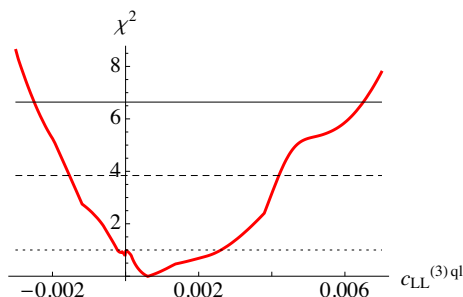


Figure 3. The χ^2 -distribution for the Wilson coefficient of the operator $\mathcal{O}_{LL}^{(3)ql}$.

in the signal plus a 4.4% in the luminosity are summed in quadrature). We checked that such estimates reproduced satisfactorily the results in the case of the HNC model. Our results, in form of the χ^2 -distribution, are summarized in figure 3 and imply

$$-0.001 \lesssim c_{LL}^{(3)ql} \lesssim 0.004 \quad \text{at 95\% C.L.} \quad . \quad (\text{B.1})$$

Combining it with the muon channel can give better limits in flavour-universal BSM models, as the ones considered here. Notice that bounds on $c_{LL}^{(3)ql}$ from the differential distribution of $pp \rightarrow l^+l^-$ at the LHC [24] are weakened once the contributions from other operators are taken into account. Therefore, eq. (B.1) provides at present the best model-independent bound on this Wilson coefficient.

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