

Higgs Physics

A. Pomarol

Dept. de Física, Universitat Autònoma de Barcelona, Barcelona, Spain

Abstract

With the discovery of the Higgs, we have access to a plethora of new physical processes that allow us to further test the SM and beyond. We show a convenient way to parametrize these physics using an effective theory for Higgs couplings, discussing the importance of the basis selection, predictions from a SM effective field theory, and possible ways to measure these couplings with special attention to the high-energy regime. Predictions from the MSSM and MCHM, with the comparison with data, are also provided.

1 Motivation

The 4th of July of 2012 marked a milestone in particle physics, as CERN announced the discovery of a new particle whose properties were in accordance with the sought-after Higgs boson [1]. Since then, we have been accumulating more and more data and measuring more decay channels, increasing the significance of the discovery while keeping at the same time a good agreement with the predictions from the Standard Model (SM) Higgs [2, 3]. To appreciate this agreement, it is convenient to plot the experimental fit to Higgs couplings in the coupling–mass plane, as shown in Fig. 1 by courtesy of CMS [2]. Were this new particle not the SM Higgs, we would have expected its couplings to lay on any point of this plane, and therefore differing significantly from the SM predictions. As an example, let us consider a scalar coming from a weak-doublet not being (the main) responsible for electroweak symmetry breaking (EWSB). This scalar could have couplings to fermions as large as $O(1)$, but very small couplings to Z/W . These predictions are shown in red in Fig. 1. Data clearly disfavours this type of scalars as compared with the SM Higgs whose predictions lay on a straight line. We can then say today that the SM Higgs is significantly supported by the experimental data, leaving most competitors far behind.

Having discovered the Higgs, we have now experimental access to new processes that will help us to test the SM and beyond. There is a fundamental aspect that makes Higgs physics very special: the Higgs is the only particle of the SM that its lightness ($m_h \sim 125 \text{ GeV} \ll M_P$) is not expected on theoretical grounds, requiring the presence of new physics beyond the SM (BSM) at the TeV. This is referred as the hierarchy problem. This makes the Higgs boson one of the most sensitive SM particle to BSM effects, and therefore the measurement of its properties one of the best ways to indirectly discover new physics and help to discriminate between different BSMs. As an example, two of the most well-motivated BSM scenarios, the minimal supersymmetric SM and the composite Higgs, predict, as we will see below, sizeable corrections to the Higgs couplings. In few words, natural theories explaining the lightness of the Higgs demand the Higgs to be SM-like only in a first approximation, predicting departures from the SM predictions to be seen in the near future.

2 Effective Higgs couplings

To characterize the most interesting Higgs processes, it is convenient to parametrize, in the most general way possible, the couplings of the Higgs to the SM particles. For this purpose we will write an effective theory for the Higgs couplings, \mathcal{L}_h . We will define \mathcal{L}_h in position-space, as it makes it simpler to eliminate redundancies. Our only approximation at this point will be to assume that the momenta q in the Higgs form-factors are smaller than a heavy scale Λ associated with the BSM physical scale, $q/\Lambda \ll 1$. This is equivalent to say that we can make an expansion in derivatives D_μ/Λ in \mathcal{L}_h . We leave for later the implications when an expansion of SM fields over Λ can be also carried out. We assume that the interactions preserve $SU(3)_c \times U(1)_{EM}$, with the Higgs defined as a neutral CP-even scalar field.

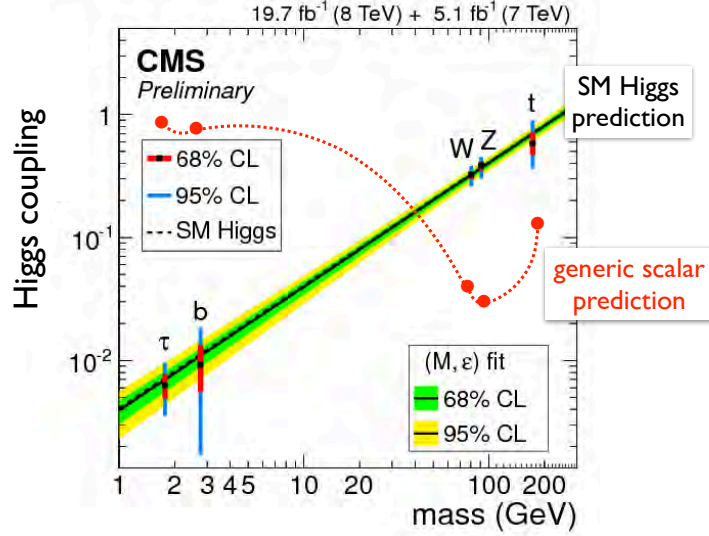


Fig. 1: Fit of the Higgs couplings, g_{ff}^h and $\sqrt{g_{VV}^h}/2v$, and predictions from the SM [2]. A generic scalar would have couplings to the SM particles laying in any point of this plane, as the example shown in red. The experimental data clearly favors a SM Higgs.

We split the Higgs couplings in two sets. One set that consists of what we call *primary* Higgs couplings and the other set containing the rest. These primaries, as we will explain later, play an important role, both theoretically and phenomenologically. We then write

$$\mathcal{L}_h = \mathcal{L}_h^{\text{primary}} + \Delta\mathcal{L}_h. \quad (1)$$

We will only keep interactions up to order $O(h^3)$, $O(h\partial^2 V^2)$ and $O(hVf^2)$ since they are the most relevant for Higgs phenomenology (adding more derivatives will be suppressed by inverse powers of Λ , and adding more fields makes the interactions harder to be observed at colliders since they will be further suppressed by phase space). Then, for CP-conserving couplings, we have *without loss of generality*¹

$$\begin{aligned} \mathcal{L}_h^{\text{primary}} &= g_{VV}^h h \left[W^{+\mu} W_{\mu}^{-} + \frac{1}{2c_{\theta_W}^2} Z^{\mu} Z_{\mu} \right] + \frac{1}{6} g_{3h} h^3 + g_{ff}^h (h\bar{f}_L f_R + h.c.) \\ &+ \kappa_{GG} \frac{h}{2v} G^{A\mu\nu} G_{\mu\nu}^A + \kappa_{\gamma\gamma} \frac{h}{2v} A^{\mu\nu} A_{\mu\nu} + \kappa_{Z\gamma} \frac{h}{v} A^{\mu\nu} Z_{\mu\nu}, \end{aligned} \quad (2)$$

and

$$\begin{aligned} \Delta\mathcal{L}_h &= \delta g_{ZZ}^h h \frac{Z^{\mu} Z_{\mu}}{2c_{\theta_W}^2} + g_{Zff}^h \frac{h}{2v} (Z_{\mu} J_N^{\mu} + h.c.) + g_{Wff'}^h \frac{h}{v} (W_{\mu}^{+} J_C^{\mu} + h.c.) \\ &+ \kappa_{WW} \frac{h}{v} W^{+\mu\nu} W_{\mu\nu}^{-} + \kappa_{ZZ} \frac{h}{2v} Z^{\mu\nu} Z_{\mu\nu}, \end{aligned} \quad (3)$$

where $J_N^{\mu} = \bar{f}\gamma^{\mu}f$ (for $f = f_L, f_R$) and $J_C^{\mu} = \bar{f}\gamma^{\mu}f'$ are respectively the neutral and charged currents. Flavour indices are implicit. We also defined $c_{\theta_W} \equiv \cos\theta_W$ where θ_W is the weak-angle, and $G_{\mu\nu}^A \equiv \partial_{\mu}G_{\nu}^A - \partial_{\nu}G_{\mu}^A$ for gluons, and similarly for the photon, A_{μ} , the Z_{μ} and W_{μ}^{+} . We can use field redefinitions to rewrite the couplings in Eq. (2) and Eq. (3) in a different way. For example, some linear combinations of the contact-interactions $hV_{\mu}J^{\mu}$ could be written as interactions of the type $hV_{\mu}\partial_{\nu}F^{\mu\nu}$ [4] by the redefinition $V_{\mu} \rightarrow (1 + \alpha h)V_{\mu}$, with an appropriate α , in the full Lagrangian (and using integration by parts). Nevertheless, we consider that Eq. (2) and Eq. (3) are the most convenient

¹From here and on, all Higgs-coupling coefficients are defined real.

way to write the Higgs couplings. Our parametrization of Higgs couplings gives priority to operators with the largest number of fields (as opposed to operators with more derivatives), as this is important when estimating the size of the couplings or looking for the dominant effects in the high-energy regime, as we will show later.

For CP-violating couplings we have

$$\mathcal{L}_h^{\text{primary}} = \delta\tilde{g}_{hff} (ih\bar{f}_L f_R + h.c.) + \tilde{\kappa}_{GG} \frac{h}{2v} G^{A\mu\nu} \tilde{G}_{\mu\nu}^A + \tilde{\kappa}_{\gamma\gamma} \frac{h}{2v} A^{\mu\nu} \tilde{A}_{\mu\nu} + \tilde{\kappa}_{Z\gamma} \frac{h}{v} A^{\mu\nu} \tilde{Z}_{\mu\nu}, \quad (4)$$

$$\begin{aligned} \Delta\mathcal{L}_h &= \tilde{g}_{Zff}^h \frac{h}{2v} (iZ_\mu J_N^\mu + h.c.) + \tilde{g}_{Wff'}^h \frac{h}{v} (iW_\mu^+ J_C^\mu + h.c.) \\ &+ \tilde{\kappa}_{WW} \frac{h}{v} W^{+\mu\nu} \tilde{W}_{\mu\nu}^- + \tilde{\kappa}_{ZZ} \frac{h}{2v} Z^{\mu\nu} \tilde{Z}_{\mu\nu}, \end{aligned} \quad (5)$$

where $\tilde{G}^{A\mu\nu} = \epsilon^{\mu\nu\rho\sigma} G_{\rho\sigma}^A/2$ and similarly for other gauge bosons.

It is important to understand the implications of global symmetries in the Higgs couplings. In particular, if the Higgs couplings are induced from BSMs that respect a custodial SU(2) symmetry [5] only weakly broken by the gauging of U(1)_Y and fermions masses, and responsible for $m_W^2 = m_Z^2 c_{\theta_W}^2$ at tree-level, we have the relations [6]²

$$\kappa_{WW} = c_{\theta_W}^2 \kappa_{ZZ} + s_{2\theta_W} \kappa_{Z\gamma} + s_{\theta_W}^2 \kappa_{\gamma\gamma}, \quad (6)$$

$$\begin{aligned} c_{\theta_W} g_{Zff}^h &= \sqrt{2} T_{3f} g_{Wff'}^h V_{\text{CKM}}^\dagger - Y_f \delta g_{ZZ}^h / m_W \quad \text{for } f = \text{up-type fermion}, \\ c_{\theta_W} g_{Zf'f'}^h &= \sqrt{2} T_{3f'} V_{\text{CKM}}^\dagger g_{Wff'}^h - Y_{f'} \delta g_{ZZ}^h / m_W \quad \text{for } f' = \text{down-type fermion}, \end{aligned} \quad (7)$$

where T_{3f} and Y_f are respectively the 3-component isospin and hypercharge of the fermion f , with $Q_f = T_{3f} + Y_f$ the electric charge, and V_{CKM} the CKM quark-mixing matrix [7]. Eq. (6) was first derived in [8]. A left-right parity P_{LR} [9] can further restrict the coefficients [6]:

$$\kappa_{Z\gamma} = \frac{c_{2\theta_W}}{s_{2\theta_W}} \kappa_{\gamma\gamma}. \quad (8)$$

Similar expressions are derived for the CP-violating counterparts.

We can also have a reduction of Higgs couplings due to dynamical reasons. For example, in BSMs with a strongly-interacting Higgs, we can neglect $\kappa_{ZZ,WW}$ in comparison with g_{VV}^h and δg_{ZZ}^h , as the formers are associated to interactions that contain more derivatives and therefore are expected to be smaller in our D_μ/Λ expansion (see later a power counting for these couplings). Also in "universal" BSMs (as those in which the BSM states only couple to SM bosons and not to fermions) we only have three relevant contact-interactions $hV_\mu J^\mu$:

$$g_{ZJ_3}^h \frac{h}{v} Z_\mu J_3^\mu, \quad g_{ZJ_Y}^h \frac{h}{v} Z_\mu J_Y^\mu, \quad g_{WJ}^h \frac{h}{v} (W_\mu^+ J_W^\mu + h.c.), \quad (9)$$

where J_3^μ , J_Y^μ and J_W^μ are respectively the 3-component isospin, hypercharge and charged SM currents [7]. Demanding also custodial invariance, we obtain

$$g_{ZJ_3}^h = \frac{g_{WJ}^h}{c_{\theta_W}}, \quad g_{ZJ_Y}^h = -\frac{\delta g_{ZZ}^h}{c_{\theta_W} m_W}, \quad (10)$$

that is equivalent to

$$g_{Wff'}^h = g_{WJ}^h V_{\text{CKM}}, \quad g_{Zff}^h = T_{3f} \frac{\sqrt{2} g_{WJ}^h}{c_{\theta_W}} - Y_f \frac{\delta g_{ZZ}^h}{c_{\theta_W} m_W}. \quad (11)$$

Eq. (11), together with Eq. (6), show that universality and custodial symmetry reduce Eq. (3) to only 3 independent Higgs couplings, that we can take to be δg_{ZZ}^h , κ_{ZZ} and g_{WJ}^h . This is in accordance with [8].

²The terms proportional to Y_f arise from the operator $\partial_\mu h Z_\nu g' B^{\mu\nu}$ that, after field redefinitions, can be rewritten as interactions in Eq. (2) and Eq. (3). One has to have this in mind when estimating the size of the coefficients.

3 The SM predictions for Higgs couplings

In the SM the Higgs sector is given by

$$\mathcal{L}_h^{\text{SM}} = |D_\mu H|^2 - (y_u \bar{Q}_L \tilde{H} u_R + y_d \bar{Q}_L H d_R + y_e \bar{L}_L H e_R + h.c.) + \mu^2 |H|^2 - \lambda |H|^4, \quad (12)$$

where the complex Higgs field H is a $\mathbf{2}_{1/2}$ of $\text{SU}(2)_L \times \text{U}(1)_Y$, $\tilde{H} = i\sigma^2 H^*$, and

$$Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}, \quad L_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}. \quad (13)$$

When the Higgs gets a vacuum expectation value (VEV), $\langle H \rangle = (0 \ v/\sqrt{2})^T$, where $v \simeq 246$ GeV, the gauge bosons W/Z and fermions get a mass proportional to their coupling to the Higgs field. Out of the 4 degrees of freedom in H , 3 corresponds to the would-be Nambu-Golstone bosons that become the longitudinal component of the W and Z , and the 4th is the Higgs particle h . In the SM all couplings of the Higgs are predicted as a function of particle masses. We have, at tree-level, that the only nonzero couplings are

$$g_{ff}^h = -\frac{gm_f}{2m_W}, \quad g_{VV}^h = gm_W, \quad g_{3h} = -\frac{3gm_h^2}{2m_W}, \quad (14)$$

that lead to the straight line of Fig. 1. The rest of the Higgs couplings arise at the loop level; κ_{GG} is mainly induced by the top loop, while $\kappa_{\gamma\gamma}$ and $\kappa_{Z\gamma}$ are generated by W and top loops, as can be found for example in [10].

4 Higgs couplings in an Effective Field Theory approach to the SM

Let us consider BSMs characterized by a mass-scale Λ much larger than the electroweak scale m_W , such that, after integrating out the BSM sector, we can make an expansion not only in derivatives D_μ over Λ , as we did in previous sections, but also an expansion of SM fields over Λ . In this way we can obtain an Effective Field Theory (EFT) made of local operators:³

$$\mathcal{L}_{\text{EFT}} = \frac{\Lambda^4}{g_*^2} \mathcal{L} \left(\frac{D_\mu}{\Lambda}, \frac{g_* H}{\Lambda}, \frac{g_* f_{L,R}}{\Lambda^{3/2}}, \frac{g F_{\mu\nu}}{\Lambda^2} \right) \simeq \mathcal{L}_4 + \mathcal{L}_6 + \dots \quad (15)$$

Here \mathcal{L}_d denotes the term in the expansion made of local operators of dimension d , while g_* denotes a generic coupling, and g and $F_{\mu\nu}$ represent respectively the SM gauge couplings and field-strengths. The Lagrangian in Eq. (15) is based on dimensional analysis and the dependence on the coupling g_* is easily obtained when the Planck constant \hbar is put back in place. Indeed, working with units $\hbar \neq 1$, the couplings have dimensions $[g_*] = [\hbar]^{-1/2}$, while $[H] = L^{-1} \cdot [\hbar]^{1/2}$ and the Lagrangian mass-terms $[\Lambda] = L^{-1}$. This dictates the dimensionless expansion-parameters to be $g_* H/\Lambda$ and D_μ/Λ , and that terms in the Lagrangian that contains n fields must carry $n - 2$ couplings to have the right dimensions. This counting is therefore valid even if g_* is not small. Although we are using a generic coupling and mass-scale, g_* and Λ , it is clear that this ought not to be always the case. For example, for a strongly-interacting light Higgs (SILH) [4] only the couplings of the Higgs to the strong BSM sector are large ($g_* \gg 1$ for the Higgs), while SM fermions are assumed to have small couplings ($g_* \sim \sqrt{y_f}$ for fermions).

The Lagrangian terms of \mathcal{L}_4 redefine the SM (and have no physical impact), while \mathcal{L}_6 encodes the dominant BSM effects. Therefore the study of the physical implications of \mathcal{L}_6 in the physics of the SM is of great importance. There are different bases used in the literature for the set of independent $d = 6$ operators in \mathcal{L}_6 . Although physics is independent of the choice of basis, it is clear that some bases are better suited than others in order to extract the relevant information, *e.g.*, for Higgs physics.

³This EFT also contains operators of dimension five, \mathcal{L}_5 , but these induce neutrino masses and therefore their coefficients must be very small (or their suppression scale Λ very large). For this reason we neglect them here since they cannot play any role for Higgs physics at the TeV.

The first complete and non-redundant basis of dimension-6 operators was given in [11]. The virtue of that basis is that it is constructed with the maximum number of operators made of fields instead of using derivatives, following our approach for Eq. (2) and Eq. (3). As we mentioned, this can be useful when estimating the size of the coefficients (see section below) or looking for the dominant effects at high-energies. Nevertheless, from a model-building point of view, it can be more advantageous to define bases that capture in few operators the impact of the most interesting BSM scenarios. With this philosophy, the SILH basis was constructed in [4], and generalised to a complete \mathcal{L}_6 basis in [8,9]. In this basis "universal" BSMs are encoded in few operators made only of SM bosons. This has the virtue of, for example, having a more direct connection between operator coefficients and the S and T parameters [12] that characterize the main electroweak effects of these BSMs. This simplicity is not present in the basis of [11] in which the equivalent of the S and T parameters involve vertex corrections [13] and then a less direct connection with the operator coefficients. Another useful basis is given in [14] with the interesting property of having a one-to-one correspondence between operators and the most relevant physical interactions measured at experiments.

In all the above mentioned bases it is possible to separate the operators into the following two groups: those that could (in principle) be induced *at tree-level* from integrating out heavy states with spin ≤ 1 in renormalizable weakly-interacting BSMs, and those operators that can only be induced *at the one-loop level* from these BSMs [9,15]. This property is, however, not respected for bases constructed with the operators of [16] where *tree* and *loop* operators are mixed.

The coefficients of \mathcal{L}_6 , referred as Wilson coefficients, are generated at the scale Λ where they are generated after integrating out the BSM heavy states. The renormalization group evolution (RGE) from Λ down to the electroweak scale, where they are supposed to be measured, can give important corrections to the Wilson coefficients and mix them [9,17,18]. For example, in supersymmetric theories or composite Higgs models, where the Wilson coefficients can be determined (see below), the RGE give us the leading-log corrections to the predictions for the Higgs couplings at low-energy that can be significant in certain cases [9].

The full set of physical implications of \mathcal{L}_6 was given in [13], where it was shown that not all type of interactions can be obtained from \mathcal{L}_6 and, of the possible ones, not all of them are independent. The set of independent couplings that are, at present, the experimentally best tested ones, were called *primary* couplings. The ones of the Higgs are presented below.

4.1 Primary Higgs couplings

Among all dimension-6 operators present in \mathcal{L}_6 , there are few of them that contribute *only* to Higgs couplings and not to other couplings (such as Vff) [9]. These are the set of independent dimension-6 operators constructed with $|H|^2$. The CP-conserving ones are ⁴

$$\begin{aligned}
 & |H|^2 \bar{Q}_L \tilde{H} u_R + h.c. , \quad |H|^2 \bar{Q}_L H d_R + h.c. , \quad |H|^2 \bar{L}_L H e_R + h.c. , \\
 & |H|^2 |D_\mu H|^2 , \quad |H|^6 , \quad |H|^2 G^{A\mu\nu} G_{\mu\nu}^A , \quad |H|^2 B^{\mu\nu} B_{\mu\nu} , \quad |H|^2 W^{a\mu\nu} W_{\mu\nu}^a ,
 \end{aligned} \tag{16}$$

where W_μ^a, B_μ are the $SU(2)_L \times U(1)_Y$ gauge bosons. To see that, indeed, the above operators can *only* be probed by measuring Higgs couplings, we just have to put the Higgs field in the EWSB vacuum, $|H|^2 \rightarrow v^2/2$, and realize that the resulting terms are operators already present in the SM, *i.e.*, their only effect is a redefinition of the SM parameters.

The set of Higgs couplings that can be independently generated from Eq. (16) are the primary Higgs couplings [13]. Their measurements provide new probes to new physics only accessible by Higgs physics. The number of primary Higgs couplings must obviously coincide with the number of Wilson coefficients associated with the operators of Eq. (16) (for the CP-conserving case). We have chosen as primary Higgs couplings those in Eq. (2), as all of them can be independently generated from the

⁴Notice that the operator $|H|^2 f \not{D} f$ can always be eliminated from the Lagrangian by field redefinitions.

operators of Eq. (16). We must be aware however that the correspondence is not one-to-one [9, 19]. There is a certain freedom to choose the set of primary Higgs couplings. For example, instead of $\kappa_{\gamma\gamma}$ and $\kappa_{Z\gamma}$, we could have taken $\kappa_{ZZ,WW}$, as these latter can also receive independent contributions from Eq. (16). The reason to choose Eq. (2) as primary Higgs couplings it is just experimental: they are the set of primary Higgs couplings best measured at the LHC.

Similarly, the CP-violating dimension-6 operators constructed with $|H|^2$ are

$$i|H|^2\bar{Q}_L\tilde{H}u_R + h.c. , \quad i|H|^2\bar{Q}_L H d_R + h.c. , \quad i|H|^2\bar{L}_L H e_R + h.c. , \\ |H|^2 G^{A\mu\nu}\tilde{G}_{\mu\nu}^A , \quad |H|^2 B^{\mu\nu}\tilde{B}_{\mu\nu} , \quad |H|^2 W^{a\mu\nu}\tilde{W}_{\mu\nu}^a , \quad (17)$$

that can independently generate the set of primary Higgs couplings of Eq. (4). Again, all these operators for $|H|^2 \rightarrow v^2/2$ generate SM terms (that redefine SM parameters) and therefore their physical effects can only be seen in Higgs physics.

The primary Higgs couplings can enter at the quantum level in other non-Higgs observables. For example, the CP-violating Higgs couplings can contribute at the loop-level to the neutron and electron electric dipole moment (EDM). The fact that we have excellent bounds on these EDMs, place indirect bounds on these Higgs couplings. We must be aware however that these bounds are model-dependent, as there can be, in principle, other BSM effects entering in the EDMs.

4.2 Beyond the primaries

The rest of CP-conserving Higgs couplings, beyond the primaries, are those of Eq. (3) at the order we mentioned before. They can in principle be generated from operators in \mathcal{L}_6 .⁵ Nevertheless, it can be proven [9, 19] that contributions from \mathcal{L}_6 to Eq. (3) are not independent from contributions to primary Higgs couplings and other electroweak couplings. Therefore they can, in principle, be constrained by other experimental measurements. As an example, consider the operator $H^\dagger D_\mu H \bar{e}_R \gamma^\mu e_R$. This gives a contribution to the Higgs coupling g_{Zff}^h , but it also contributes to the coupling $Z\bar{e}_R e_R$ that has been very-well measured at LEP, putting strong bounds on possible BSM effects.

The explicit relations between the \mathcal{L}_6 -contributions to Eq. (3) and to other couplings were explicitly calculated in [13, 14, 19] assuming family universality. Here we give these relations for the general case (derived at the tree-level) [6]:

$$\delta g_{ZZ}^h = 2gt_{\theta_W}^2 m_W (c_{\theta_W}^2 \delta g_1^Z - \delta\kappa_\gamma) , \\ g_{Zff}^h = 2\delta g_{ff}^Z - 2\delta g_1^Z (g_{ff}^Z c_{2\theta_W} + g_{ff}^\gamma s_{2\theta_W}) + 2\delta\kappa_\gamma Y_f \frac{e s_{\theta_W}}{c_{\theta_W}^3} , \quad g_{Wff'}^h = 2\delta g_{ff'}^W - 2\delta g_1^Z g_{ff'}^W c_{\theta_W}^2 , \quad (18) \\ \kappa_{ZZ} = \frac{1}{c_{\theta_W}^2} \delta\kappa_\gamma + 2 \frac{c_{2\theta_W}}{s_{2\theta_W}} \kappa_{Z\gamma} + \kappa_{\gamma\gamma} , \quad \kappa_{WW} = \delta\kappa_\gamma + \kappa_{Z\gamma} + \kappa_{\gamma\gamma} , \quad (19)$$

with

$$\delta g_{ff'}^W = \frac{c_{\theta_W}}{\sqrt{2}} (\delta g_{ff}^Z V_{\text{CKM}} - V_{\text{CKM}} \delta g_{ff'}^Z) \quad \text{for } f = f_L , \quad (20)$$

and where

$$g_{ff}^\gamma = eQ_f , \quad g_{ff}^Z = \frac{g}{c_{\theta_W}} (T_{3f} - Q_f s_{\theta_W}^2) , \quad g_{ff'}^W = \frac{g}{\sqrt{2}} V_{\text{CKM}} , \quad 0 \text{ resp. for } f = f_L, f_R , \quad (21)$$

are the γ , Z and W couplings to fermions in the SM. Flavor indices are again implicit. We have also defined by δg_{ff}^Z ($\delta g_{ff'}^W$) the BSM corrections to the Z (W) couplings to fermions:

$$\Delta\mathcal{L}_{ff}^V = \frac{\delta g_{ff}^Z}{2} (Z_\mu J_N^\mu + h.c.) + \delta g_{ff'}^W (W_\mu^+ J_C^\mu + h.c.) , \quad (22)$$

⁵At $O(hFff)$ we also have dipole-type interactions that can arise from \mathcal{L}_6 . Their Wilson coefficients are however expected to be suppressed by SM Yukawa-couplings (otherwise could largely contribute at the loop level to the SM fermion masses). These couplings are related to fermion EDMs as can be found in [13].

while δg_1^Z is the correction to the ZWW coupling and $\delta\kappa_\gamma$ parametrizes BSM contributions to the EDM of the W , following the notation of [16] for anomalous triple gauge couplings (TGC):

$$\Delta\mathcal{L}_{3V} = igc_{\theta_W}\delta g_1^Z \left[Z^\mu (W^{+\nu}W_{\mu\nu}^- - h.c.) + Z^{\mu\nu}W_\mu^+W_\nu^- \right] + ie\delta\kappa_\gamma \left[(A^{\mu\nu} - t_{\theta_W}Z^{\mu\nu})W_\mu^+W_\nu^- \right]. \quad (23)$$

Following [13], we have chosen to work in the mass-eigenstate basis within a parametrization in which kinetic terms and masses do not receive corrections and then take the SM values. All BSM effects are in couplings. We think this is the most convenient parametrization of BSM effects due to the straightforward connection between couplings and physical processes, that in most of the cases is a one-to-one correspondence. The SM input parameters can be taken to be α_{EM} , m_Z and m_W that, in our parametrization, do not have BSM corrections, as opposed to G_F that receive corrections from 4-fermion interactions. We remark again that the predictions Eq. (18) and Eq. (19) are derived at the tree-level and only apply to BSM effects coming from \mathcal{L}_6 . There are also SM contributions to these couplings at the loop level, that can be as important as new-physics contributions, and must be incorporated accordingly.

Eq. (18) and Eq. (19) are important results. They show that all Higgs couplings of Eq. (3) can be written as a function of BSM effects to two primary Higgs couplings ($\kappa_{\gamma\gamma}$, $\kappa_{Z\gamma}$), Z/W couplings to SM fermions (δg_{ff}^Z , $\delta g_{fRf'R}^W$), and two TGC (δg_1^Z , $\delta\kappa_\gamma$). Experimental bounds on $\kappa_{\gamma\gamma, Z\gamma}$ are already at the per-cent level [19], while Z/W couplings have also been experimentally constrained, mostly from LEP and SLC [20, 21] (with Tevatron providing an accurate measurement of the W -mass). One finds that bounds on δg_{ff}^Z are quite strong, at the per mille-level in most of the cases, but bounds on δg_1^Z and $\delta\kappa_\gamma$ are much weaker [22]. Therefore, at present, we can already derive, using Eq. (18) and Eq. (19), relevant model-independent bounds on the Higgs couplings of Eq. (3) [19].

In the case of custodial-invariant universal BSMs, Eq. (18) reduces to

$$\begin{aligned} \delta g_{ZZ}^h &= 2gt_{\theta_W}^2 m_W \left(c_{\theta_W}^2 \delta g_1^Z - \delta\kappa_\gamma + \widehat{S} \right), \\ g_{Zff}^h &= -2T_{3f} \delta g_1^Z g_{c_{\theta_W}} - Y_f \frac{\delta g_{ZZ}^h}{c_{\theta_W} m_W}, \quad g_{Wff'}^h = -2\delta g_1^Z g_{ff'}^W c_{\theta_W}^2, \end{aligned} \quad (24)$$

where \widehat{S} is, up to a normalization constant [23], the S -parameter [12]. As expected, Eq. (24) and Eq. (18) fulfill Eq. (6) and Eq. (11), and g_{Zff}^h is fully determined by the custodial symmetry as a function of $g_{Wff'}^h$ and δg_{ZZ}^h .

The CP-violating non-primary Higgs couplings, Eq. (5), are also not independent but related to other couplings. We have

$$\begin{aligned} \tilde{g}_{Zff}^h &= 2\delta\tilde{g}_{ff}^Z, & \tilde{g}_{Wff'}^h &= 2\delta\tilde{g}_{ff'}^W, \\ \tilde{\kappa}_{ZZ} &= \frac{1}{c_{\theta_W}^2} \delta\tilde{\kappa}_\gamma + 2\frac{c_{2\theta_W}}{s_{2\theta_W}} \tilde{\kappa}_{Z\gamma} + \tilde{\kappa}_{\gamma\gamma}, & \tilde{\kappa}_{WW} &= \delta\tilde{\kappa}_\gamma + \tilde{\kappa}_{Z\gamma} + \tilde{\kappa}_{\gamma\gamma}, \end{aligned} \quad (25)$$

where

$$\delta\tilde{g}_{ff'}^W = \frac{c_{\theta_W}}{\sqrt{2}} (\delta\tilde{g}_{ff}^Z V_{\text{CKM}} - V_{\text{CKM}} \delta\tilde{g}_{f'f'}^Z) \text{ for } f = f_L, \quad (26)$$

with $\delta\tilde{g}_{ff}^Z$ and $\delta\tilde{g}_{ff'}^W$ defined as

$$\Delta\tilde{\mathcal{L}}_{ff}^V = \frac{\delta\tilde{g}_{ff}^Z}{2} (iZ_\mu J_N^\mu - h.c.) + \delta\tilde{g}_{ff'}^W (iW_\mu^+ J_C^\mu - h.c.), \quad (27)$$

and $\tilde{\kappa}_\gamma$ being the CP-violating TGC:

$$\Delta\mathcal{L}_{3\tilde{V}} = ie\delta\tilde{\kappa}_\gamma \left[(\tilde{A}^{\mu\nu} - t_{\theta_W} \tilde{Z}^{\mu\nu}) W_\mu^+ W_\nu^- \right]. \quad (28)$$

The predictions Eq. (18), Eq. (19) and Eq. (25) rely on the (quite plausible) hypothesis that the leading SM deviations arise from \mathcal{L}_6 . Finding experimental evidence for deviations from these predictions, would mean that nature does not fulfil this hypothesis: either because there are light BSM states ($\Lambda \lesssim m_h$), the composite-scale of the Higgs is low ($\Lambda \sim g_* v$), that is equivalent to say that h cannot be identified within the SM doublet, or that there are other sources of EWSB independent of $\langle H \rangle$ [24].

4.3 Power counting for Higgs couplings

It can be useful to estimate the size of the contributions to the effective Higgs couplings arising from generic BSMs. As it is clear from the expansion in Eq. (15), the coefficients in Eq. (2) and Eq. (3) can have different dependence with g_* . The Higgs couplings that can receive the largest power of g_* are g_{3h} and g_{ff}^h where

$$\delta g_{3h} \sim \frac{g_*^4 v^3}{\Lambda^2}, \quad \delta g_{ff}^h \sim \frac{g_*^3 v^2}{\Lambda^2}. \quad (29)$$

For $g_* \gg 1$, Eq. (29) can give $O(1)$ corrections to g_{3h} and g_{ff}^h , even after demanding $g_*^2 v^2 / \Lambda^2 \ll 1$ necessary to make the expansion Eq. (15) valid. Nevertheless, in theories where the Higgs mass is protected by a symmetry, as it happens in theories that solve the hierarchy problem such as supersymmetry or composite Higgs models, the contributions to g_{3h} are also expected to be protected and then proportional to $m_h^2 / v^2 \sim \lambda$. Also it is natural to expect that chirality protects terms proportional to $\bar{f}_L f_R$, at least by a Yukawa coupling $y_f \sim m_f / v$, otherwise corrections to fermion masses would be too large. For this reason, it is more natural to assume that the corrections to these Higgs couplings are of order

$$\delta g_{3h} \sim \lambda v \frac{g_*^2 v^2}{\Lambda^2}, \quad \delta g_{ff}^h \sim y_f \frac{g_*^2 v^2}{\Lambda^2}, \quad (30)$$

that potentially give relative corrections of $O(g_*^2 v^2 / \Lambda^2)$. At the same order, we also have

$$\delta g_{VV}^h \sim g^2 v \frac{g_*^2 v^2}{\Lambda^2}, \quad (31)$$

and

$$\delta g_{ff}^Z, \delta g_1^Z \sim g \frac{g_*^2 v^2}{\Lambda^2}. \quad (32)$$

Finally, couplings coming from a derivative (or field-strength) expansion, the κ_i , are expected to scale as

$$\kappa_i \sim \frac{g^2 v^2}{\Lambda^2}. \quad (33)$$

Nevertheless, in renormalizable BSMs these coefficients can only be induced at the loop-level and therefore expected to be

$$\kappa_i \sim \frac{g_*^2}{16\pi^2} \frac{g^2 v^2}{\Lambda^2}. \quad (34)$$

Indeed, it can be shown [4, 9] that the κ_i cannot be generated *at tree-level* from integrating out scalars, fermions and vector bosons in renormalizable theories.

The above estimates are useful to determine which are the most sizeable BSM corrections to the Higgs couplings. For example, in theories in which the Higgs is strongly coupled, the largest corrections are those of Eq. (30) and Eq. (31) that depend quadratically in the strong coupling $g_* \gg 1$ [4]. If also the SM fermions are strongly-coupled, Eq. (32) can also give similar size corrections. It is also important to remark that even for theories in which the field expansion in Eq. (15) is not valid (*e.g.*, when $g_* v \sim \Lambda$), the power counting for Higgs couplings given here is expected to be correct. In particular, the above estimates are in accordance with the NDA analysis of [25] proposed for QCD.

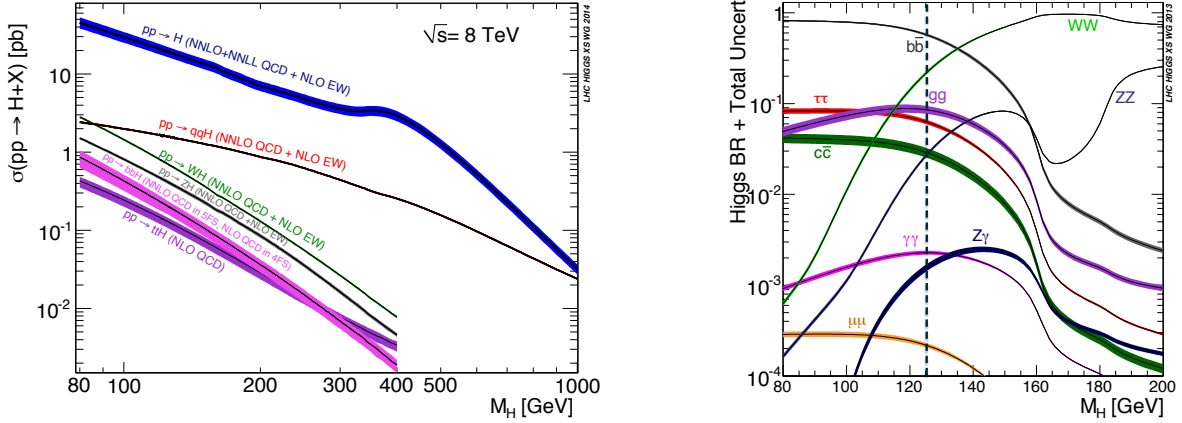


Fig. 2: Predictions for the main Higgs production cross-sections and Higgs BR in the SM [26].

For the non-primary Higgs couplings we have the estimates

$$\frac{\delta g_{ZZ}^h}{g^2 v}, \frac{g_{Zff}^h}{g}, \frac{g_{Wff'}^h}{g} \sim \frac{g_*^2 v^2}{\Lambda^2} \quad \text{and} \quad \kappa_{ZZ}, \kappa_{WW} \sim \frac{g^2 v^2}{\Lambda^2}, \quad (35)$$

in agreement with the relations in Eq. (18). Similar estimates follow for CP-violating Higgs couplings.

5 Experimental determination of the effective Higgs couplings

The primary Higgs couplings can be determined by searching for the Higgs through the different production mechanisms and decays. The main Higgs production mechanisms at the LHC are

$$\begin{aligned} \text{Gluon fusion:} & \quad GG \rightarrow h, \\ Vh\text{-associated production:} & \quad q\bar{q} \rightarrow Vh, \\ \text{Vector boson fusion (VBF):} & \quad qq \rightarrow qqVV^* \rightarrow qqh, \\ htt\text{-associated production:} & \quad GG \rightarrow t\bar{t}h, \end{aligned} \quad (36)$$

while the most important Higgs branching ratios (BR) are

$$BR(h \rightarrow b\bar{b}), BR(h \rightarrow \tau\bar{\tau}), BR(h \rightarrow Vf\bar{f}), BR(h \rightarrow \gamma\gamma), BR(h \rightarrow Z\gamma). \quad (37)$$

The predictions for a SM Higgs are given in Fig. 2. The Higgs mass can be mainly determined from the Higgs decay to $\gamma\gamma$ and Zff that allows to obtain

$$\begin{aligned} m_h &= 125.03^{+0.26}_{-0.27} \text{ (stat.) } ^{+0.13}_{-0.15} \text{ (syst.) GeV} \quad \text{from CMS}, \\ m_h &= 125.36 \pm 0.37 \text{ (stat.) } \pm 0.18 \text{ (syst.) GeV} \quad \text{from ATLAS}. \end{aligned} \quad (38)$$

At the LHC one can combine the different Higgs production mechanisms and BR of Eq. (36) and Eq. (37) to determine 7 primary Higgs couplings: g_{ff}^h ($f = t, b, \tau$), g_{VV}^h , κ_{GG} , $\kappa_{\gamma\gamma}$ and $\kappa_{Z\gamma}$.⁶ The CMS fit of six of the primary Higgs couplings is shown in Fig. 3, where other Higgs couplings have been set to zero.⁷ The fit shows a good agreement with the SM predictions and no sign of new-physics. The implications of these measurements in particular BSMs will be discussed in the next section. The

⁶We note that g_{tt}^h and g_{VV}^h also affect $BR(h \rightarrow \gamma\gamma/Z\gamma)$ and $\sigma(GG \rightarrow h)$ at the one-loop level [4].

⁷The ATLAS results are not shown here since the fit is performed only for few primaries at each time instead of a global fit to all of them [3]. For a combination of ATLAS and CMS data see, for example, [28, 29].

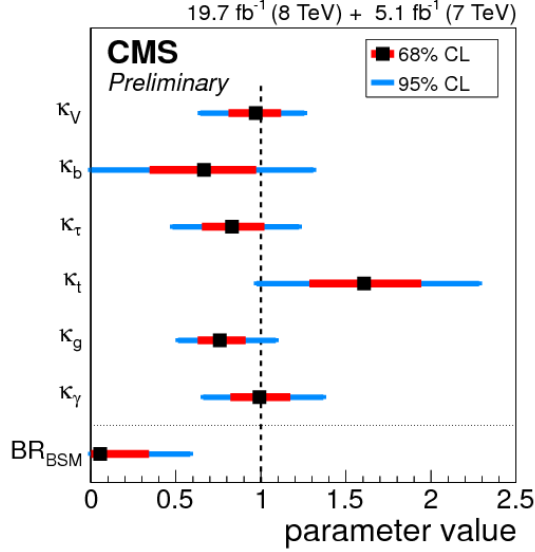


Fig. 3: Fit of 6 primary Higgs couplings from CMS [2]. Notation: $\kappa_V \equiv g_{VV}^h/g_{VV}^{hSM}$, $\kappa_f \equiv g_{ff}^h/g_{ff}^{hSM}$, $\kappa_g \equiv \kappa_{GG}/\kappa_{GG}^{SM}$ and $\kappa_\gamma \equiv \kappa_{\gamma\gamma}/\kappa_{\gamma\gamma}^{SM}$; loop effects in κ_{GG} and $\kappa_{\gamma\gamma}$ are not included [2].

primary coupling $\kappa_{Z\gamma}$ has not been included in the fit of Fig. 3, but one can use the experimental bound $BR(h \rightarrow Z\gamma)/BR(h \rightarrow Z\gamma)_{SM} \lesssim 10$ [27] to derive the constraint $-0.01 \lesssim \kappa_{Z\gamma} \lesssim 0.02$ [19]. The fact that in the SM $h \rightarrow Z\gamma$ arises at the one-loop level, and therefore has a small branching fraction $BR(h \rightarrow Z\gamma) \sim 0.15\%$, makes this BR very sensitive to new-physics; it probably provides the last chance to find large BSM effects in SM Higgs couplings.

Among the remaining primary Higgs couplings to be measured we have g_{3h} . Its determination however will be very difficult since it requires to search for double-Higgs production $pp \rightarrow hh$ that has small rates [30]. Also Higgs couplings to light fermions g_{ff}^h (beyond the 3rd family) are going to be difficult to measure since we expect these couplings to be proportional to m_f/m_W (see Eq. (14) and Eq. (30)), giving then very small BR. For example, for the case of the muon, that is probably the most accessible, we have in the SM $BR(h \rightarrow \mu\mu) \sim 0.02\%$. Therefore a high luminosity at the LHC run 2 will be needed to measure this coupling. Flavour-violating Higgs couplings in g_{ff}^h can also be accessible through Higgs decays. This is particularly interesting for theories of flavour in which Yukawas are generated from the mixing of the SM fermions with heavy BSM states. The strength of these mixings are expected to be $\sim \sqrt{m_{f_i}/v}$, and therefore predicting $g_{f_i f_j}^h \sim \sqrt{m_{f_i} m_{f_j}}/v$ that can lead to sizeable flavour-violating Higgs decays. In particular, one has $BR(h \rightarrow \tau\mu) \sim m_\mu/m_\tau \times BR(h \rightarrow \tau\tau) \sim 0.4\%$ that is quite close to the present experimental bound $BR(h \rightarrow \tau\mu) < 1.57\%$ [31]. Finally, most of the CP-violating Higgs couplings are poorly measured since they appear quadratically in production rates and BR since the interference terms with the SM contributions vanish.⁸ Kinematical differential distributions can be used to measure these couplings [32], and alternative methods have been recently proposed in [33]. Nevertheless, indirect bounds on most of these couplings are very strong (see for example [34] for bounds on $\tilde{\kappa}_{\gamma\gamma}$ from EDMs), making difficult to believe that Higgs CP-violating couplings are sizeable. The exception is probably $\delta\tilde{g}_{\tau\tau}^h$ whose bounds are not so strong and could have possible impact in CP-violating Higgs decays.

The experimental full extraction of all Higgs couplings, including the non-primary ones, Eq. (3) and Eq. (5), is a difficult task. The best way to disentangle the effects of δg_{ZZ}^h , $\kappa_{ZZ,WW}$ and g_{Vff}^h ($V = Z, W$), as well as their CP-violating counterparts, is by looking for modifications in differential

⁸Since in the SM the hGG , $h\gamma\gamma$ and $hZ\gamma$ couplings are small (as they arise at the one-loop level), the interference terms are also small, and the corresponding bounds on CP-conserving and CP-violating couplings, κ_i and $\tilde{\kappa}_i$, are comparable.

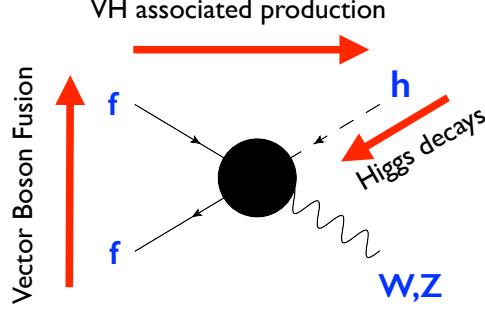


Fig. 4: The form-factor $hVff$, that as a function of the effective Higgs couplings is given in Eq. (39), can be tested in three different Higgs processes at the LHC: either in Higgs decays $h \rightarrow Vff$, in Vh -associated production or in the VBF-like process $pp \rightarrow qqV/qqVV^* \rightarrow qqh$.

distributions in Higgs processes. The most relevant ones are the Higgs decays $h \rightarrow Vff$, the Vh -associated production and the VBF-like process $pp \rightarrow qqV/qqVV^* \rightarrow qqh$. All of them arise from the $hVff$ amplitude (see Fig. 4) given by (neglecting fermion masses)

$$\mathcal{M}_{hVff}(q, p) = \frac{1}{v} \epsilon^{*\mu}(q) J_V^\nu(p) [A^V \eta_{\mu\nu} + B^V (p \cdot q \eta_{\mu\nu} - p_\mu q_\nu) + C^V \epsilon_{\mu\nu\rho\sigma} p^\rho q^\sigma], \quad (39)$$

where q and p are respectively the total 4-momentum of V and the fermion pair in $J_V^\mu = J_N^\mu, J_C^\mu$ for $V = Z, W$, and ϵ^μ is the polarization 4-vector of V . We have defined

$$A^V = a^V + \hat{a}^V \frac{m_V^2}{p^2 - m_V^2}, \quad B^V = b^V \frac{1}{p^2 - m_V^2} + \hat{b}^V \frac{1}{p^2}, \quad C^V = c^V \frac{1}{p^2 - m_V^2} + \hat{c}^V \frac{1}{p^2}, \quad (40)$$

with $\hat{b}^W, \hat{c}^W = 0$, and where

$$\begin{aligned} a^Z &= \delta g_{Zff}^h + i\delta \tilde{g}_{Zff}^h, & a^W &= \delta g_{Wff'}^h + i\delta \tilde{g}_{Wff'}^h, \\ \hat{a}^Z &= 2g_{ff}^Z \left(1 + \frac{\delta g_{VV}^h + \delta g_{ZZ}^h}{gm_W} \right), & \hat{a}^W &= 2g_{ff'}^W \left(1 + \frac{\delta g_{VV}^h}{gm_W} \right), \\ b^Z &= -2g_{ff}^Z \kappa_{ZZ}, & b^W &= -2g_{ff'}^W \kappa_{WW}, \\ \hat{b}^Z &= -2eQ_f \kappa_{Z\gamma}, & & \\ c^Z &= -2g_{ff}^Z \tilde{\kappa}_{ZZ}, & c^W &= -2g_{ff'}^W \tilde{\kappa}_{WW}, \\ \hat{c}^Z &= -2eQ_f \tilde{\kappa}_{Z\gamma}. & & \end{aligned} \quad (41)$$

From the differential distributions of the decay products in $h \rightarrow Vff$, one can put bounds on the coefficients of Eq. (40) and, consequently, on non-primary Higgs couplings. Nevertheless, we still have poor statistics and bounds on Higgs couplings are almost irrelevant unless we turn on one by one [32]. At present, the most promising way to obtain significant bounds in some of the Higgs couplings of Eq. (3) is, as we will discuss below, by measuring them at the LHC high-energy regime, for example in the Vh -associated Higgs production where the effects of some of these couplings are enhanced.

Since primary Higgs couplings predict equal deviations in the $hZff$ and $hWff$ physical amplitudes (normalized to their SM values), measuring a relative deviation between these two would provide evidence for non-primary Higgs couplings. At the LHC this relative deviation is parametrized by $\lambda_{WZ} - 1$ [2, 3] that at present does not show any evidence of being different from zero; from the experimental data we have $-0.35 < \lambda_{WZ} - 1 < 0.08$ [3]. This quantity is predicted in the SM EFT of Eq. (15) to be [19]

$$\lambda_{WZ}^2 - 1 \simeq 0.6\delta g_1^Z - 0.5\delta \kappa_\gamma - 0.7\kappa_{Z\gamma}, \quad (42)$$

where we have used Eqs. (18)-(19), neglecting $\kappa_{\gamma\gamma}$ and $\delta g_{ff}^{Z,W}$, since they are experimentally constrained to be less than $10^{-2} - 10^{-3}$.

5.1 Towards the high-energy regime

One of the most interesting perspectives at the LHC run 2 is the access to physical processes at much higher energies. This can be used to probe Higgs production mechanism or off-shell Higgs mediated processes in a regime in which the effects of some anomalous Higgs couplings can be enhanced by factors E^2/Λ^2 . As an example, let us consider the associated Higgs production, $pp \rightarrow Vh$. As it is clear from Eq. (40), at high-energies, $E \gg m_V$, the coefficient a^V dominates the amplitude. Thanks to our parametrization for Higgs couplings, this coefficient is in one-to-one correspondence with the contact-interaction g_{Vff}^h . Indeed, at the partonic level, we have

$$\sigma(qq \rightarrow hV) \Big|_{\hat{s} \gg m_h^2} = \sigma(qq \rightarrow hV_L)_{\text{SM}} \left(1 + \frac{g_{Vff}^h}{g_{ff}^V} \frac{\hat{s}}{m_V^2} + \dots \right). \quad (43)$$

By looking at high invariant-masses for hV , it is possible to put important bounds on g_{Vff}^h [35, 36]. Nevertheless, one has to be careful that one is not probing these couplings at energies above Λ where an expansion in \hat{s}/Λ^2 would not be valid. To address this issue, the power-counting of section 4.3 is crucial. Using Eq. (35), we can write $g_{Vff}^h \equiv g c_{Vff}^h g_*^2 v^2 / \Lambda^2$ where c_{Vff}^h is a coefficient $O(1)$. Now, experimentally, due to the lack of experimental accuracy in the measurement of $pp \rightarrow Vh$ at the LHC, we can only bound at present high-energy deviations from the SM to be less than $O(1)$ [35, 36], that is equivalent to say, using Eq. (43),

$$\frac{g_{Vff}^h}{g_{ff}^V} \frac{\hat{s}}{m_V^2} < O(1) \rightarrow c_{Vff}^h \lesssim \frac{\Lambda^2}{\hat{s}} \frac{g^2}{g_*^2}. \quad (44)$$

To guarantee the validity of the expansion in \mathcal{L}_h , we must stay in the regime $\Lambda^2/\hat{s} \gg 1$. Therefore the experimental bound Eq. (44) can only be restrictive (and useful) for strongly-interacting BSMs in which $g_* \gg g$. In these scenarios we can safely use the hV -production high-energy data to obtain bounds on g_{Vff}^h at the per-cent level [36]. In models in which, in addition, the expansion of Eq. (15) is valid, bounds on g_{Vff}^h can be translated into bounds on δg_1^Z . Indeed, we have from Eq. (18), after neglecting δg_{ff}^Z due to the strong constraints from LEP, and neglecting $\delta \kappa_\gamma$ since this does not grow with g_*^2 [13],

$$\delta g_1^Z \simeq - \frac{g c_{Zff}^h}{2(g_{ff}^Z c_{2\theta_W} + e Q_f s_{2\theta_W})} \frac{g_*^2 v^2}{\Lambda^2} \simeq - \frac{g c_{Wff'}^h}{2g_{ff'}^W c_{\theta_W}^2} \frac{g_*^2 v^2}{\Lambda^2}. \quad (45)$$

From the experimental data at the high-energy regime of the hV -associated production we obtain [36]

$$-0.01 < \delta g_1^Z < 0.04 \quad (95\% \text{ CL}). \quad (46)$$

This is as competitive as the one obtained from anomalous TGC at LEP [21] and at the LHC [37].

5.2 Invisible Higgs decay

We have assumed so far that there are no more light particles than those of the SM. If there were new light states to which the Higgs could decay to, all the Higgs BRs would be reduced, changing the fit of the Higgs couplings [38]. There are well-motivated BSMs where the Higgs can decay invisibly. An example is given in [39] where the Higgs can decay to a gravitino and neutrino that interact so weakly that escape from detection. Also in certain models the Higgs can decay to dark matter that, being stable and EM neutral, also escape from detection.⁹ There are direct searches for Higgs decaying invisibly based on looking for missing energy plus a $Z/W/\gamma/jet$. The CMS bound is given in Fig. 3.

⁹Alternative effects from new light physics can be found in [40].

6 Predictions for the Higgs couplings from BSM solutions to the hierarchy problem

The simplicity of the SM Higgs-mechanism is at odds with its quantum stability. The fact that the Higgs is a scalar, a spin zero state, makes it difficult to keep it light ($m_h \ll M_P$). This problematic can be easily understood just by looking at the degrees of freedom (DOF) of a massless and massive state of spin 0, and compare them with those of a state of spin 1/2, 1, or higher. Indeed, a massless vector, as the photon, has two polarizations (2 DOF), while a massive vector has 3 polarizations. The $2 \neq 3$ guarantees that a massless vector can never get a mass by continuous variations of parameters (or quantum fluctuations); only a discrete change in the theory, increasing the DOF, can make vector massive. Similarly for fermions, we have that a charged massless fermion has 2 DOF, while a massive one has the double (left- and right-handed states), and therefore, for the same reason, massless fermions are safe from getting masses under fluctuations.¹⁰ Now, massless scalars have the same DOF as massive scalars: 1 DOF for neutral ones. Even if we start with a massless scalar at tree-level, it is not guaranteed that quantum corrections will not give it a mass.

A possible solution to keep the Higgs stable from getting a large mass is to upgrade the SM to include a symmetry relating the Higgs, a scalar, to a fermion whose mass can be stable, as we explained above. This is the case of supersymmetry. An alternative option is to assume that the Higgs is not an elementary state but a state made of elementary fermions, as pions in QCD. In this case, the Higgs arises as a composite state from a new strong-sector at the TeV. It is interesting to point out that both scenarios predicted a light Higgs. While in minimal supersymmetric versions of the SM (MSSM) the lightest-Higgs mass was expected to be in the range $m_h \lesssim 135$ GeV [41], minimal versions of composite Higgs (MCHM) predicted $115 \text{ GeV} \lesssim m_h \lesssim 185 \text{ GeV}$ [42]. The connection between the Higgs mass and the mass spectrum of resonances is of crucial phenomenological interest, since allows to obtain predictions, from the present experimental value $m_h \simeq 125$ GeV, for the heavy spectrum, either stops for the MSSM [43] or fermionic resonances for the MCHM [44, 45].

In the following, we will centre in the predictions of these models to Higgs couplings. As we emphasized in the introduction, the Higgs is usually the SM particle whose couplings are most sensitive to BSM corrections. Indeed, as we will see below, in supersymmetric theories Higgs couplings can be affected at tree-level [46], while other SM couplings are affected at the loop level. Similarly, in strongly-interacting theories in which the Higgs is composite, effects on Higgs couplings can be enhanced by a factor g_*^2 [4], that can be as large as $\sim 16\pi^2$, with respect to effects in other couplings. It is also important to remark that in BSMs trying to solve the hierarchy problem the main BSM effects in Higgs physics are captured by the primary Higgs couplings, as contributions to non-primary Higgs couplings are usually negligible. This shows once more the importance of the primaries.

6.1 The Minimal Supersymmetric SM (MSSM)

We will work in the limit in which the supersymmetric spectrum is heavier than m_h . This covers *most* of the parameter space of the MSSM, after LHC searches have pushed the superpartner masses towards the TeV regime, and none deviation from the SM has been observed. Also to accommodate $m_h \simeq 125$ GeV requires large stop masses in the MSSM [43].

The only tree-level corrections to the lightest-Higgs couplings come from the extra heavy Higgs doublet of the MSSM H' . This is due to the R -parity of the MSSM that only allow R -even field tree-level corrections. At order $v^2/M_{H'}^2$ (*i.e.*, keeping only $1/\Lambda^2$ -suppressed effects where now $\Lambda = M_{H'}$), only the Higgs couplings to fermions are affected, since corrections to hVV appear at order $v^4/M_{H'}^4$, as can be easily understood from Feynman diagrams –see Fig. 5. Deviations from the SM values for the hff

¹⁰If a fermion has no charge, it can get a Majorana-type mass without increasing the DOF, as probably is the case for the SM neutrinos. For this reason, to keep naturally massless fermions, we must assume that the fermion has some type of charge.

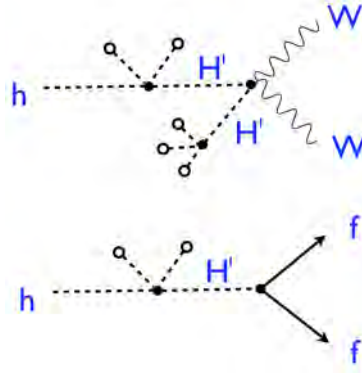


Fig. 5: Feynman diagrams contributing to g_{VV}^h and g_{ff}^h from integrating the heavy MSSM Higgs H' .

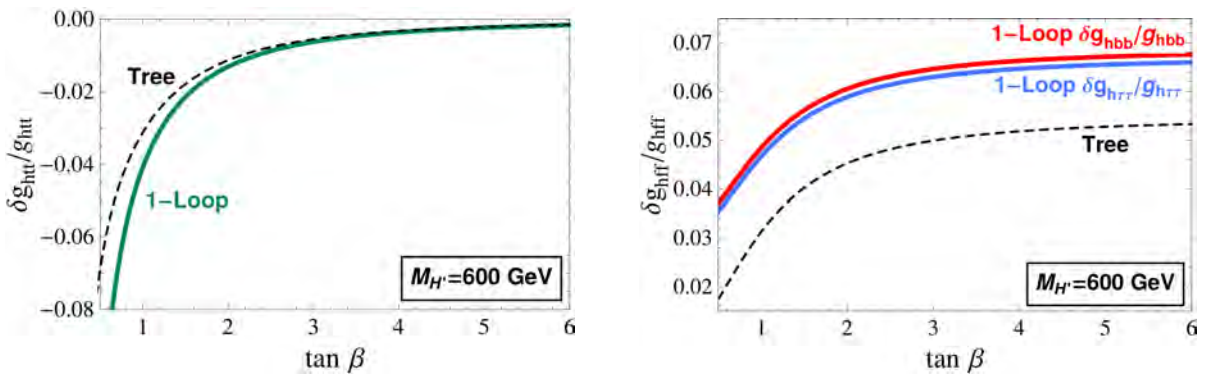


Fig. 6: Relative modifications of the Higgs couplings to fermions with respect to their SM values at tree-level (dashed line), and after including RGE effects from Λ to the electroweak scale (solid lines), as a function of $\tan \beta$ in a MSSM scenario with $\Lambda = M_{H'} = 600$ GeV and unmixed stops heavy enough to reproduce $m_h = 125$ GeV. Left plot: Higgs coupling to tops. Right plot: Higgs coupling to bottoms (upper solid line) and taus (lower solid line) [9].

couplings, including also one-loop RGE effects coming from the top, are given by [9]

$$\begin{aligned}
\frac{\delta g_{tt}^h}{g_{tt}^{h\text{SM}}} &= \frac{v^2}{M_{H'}^2} \left(\frac{\lambda'}{t_\beta} \left[1 - \frac{21y_t^2}{16\pi^2} \log \frac{M_{H'}}{m_h} \right] + \frac{3y_t^4}{4\pi^2 t_\beta^2} \log \frac{M_{H'}}{m_h} \right), \\
\frac{\delta g_{bb}^h}{g_{bb}^{h\text{SM}}} &= -\frac{v^2}{M_{H'}^2} \left(\lambda' t_\beta \left[1 - \frac{y_t^2}{2\pi^2} \log \frac{M_{H'}}{m_h} \right] + \frac{y_t^2}{16\pi^2} \left[5 \frac{\lambda'}{t_\beta} - 14y_t^2 \right] \log \frac{M_{H'}}{m_h} \right), \\
\frac{\delta g_{\tau\tau}^h}{g_{\tau\tau}^{h\text{SM}}} &= -\frac{v^2}{M_{H'}^2} \left(\lambda' t_\beta \left[1 - \frac{3y_t^2}{8\pi^2} \log \frac{M_{H'}}{m_h} \right] + \frac{3y_t^2}{8\pi^2} \left[\frac{\lambda'}{t_\beta} - 2y_t^2 \right] \log \frac{M_{H'}}{m_h} \right), \quad (47)
\end{aligned}$$

with $t_\beta \equiv \tan \beta$ and ¹¹

$$\lambda' = \frac{1}{8}(g^2 + g'^2) \sin 4\beta - \frac{3y_t^4}{8\pi^2 t_\beta} \log \frac{M_{\tilde{t}}^2}{M_{H'}^2}, \quad (48)$$

where $M_{\tilde{t}}$ is the value of the stop masses taking, for simplicity, zero stop left-right mixing. To illustrate the impact of these corrections, let us take $M_{\tilde{t}}$ large enough to get $m_h \simeq 125$ GeV through the well-known loop corrections to the Higgs quartic coupling, which at one-loop and neglecting stop left-right mixings read: $\lambda = \frac{1}{8}(g^2 + g'^2) \cos^2 2\beta + \frac{3y_t^4}{16\pi^2} \log \frac{M_{\tilde{t}}^2}{M_t^2}$. This gives the value of λ' as a function of t_β

¹¹In Eq. (48) we are also including RGE effects from $M_{\tilde{t}}$ to $M_{H'}$ proportional to the top-Yukawa y_t .

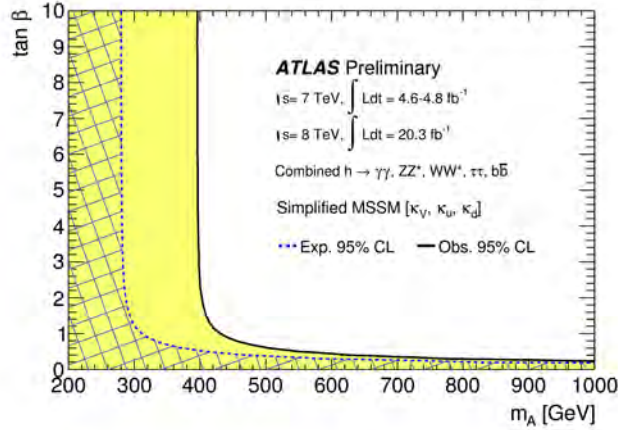


Fig. 7: Regions of the $m_A - \tan \beta$ plane excluded by Higgs physics in a MSSM with heavy partners [47].

and $M_{H'}$ that we can then plug in Eq. (47) to obtain the RGE-improved corrections for g_{ff}^h induced by integrating out the heavy Higgs. The results are shown in Fig. 6. The experimental bounds on the Higgs couplings can be translated into a bound on $M_{H'}$ as a function of $t\beta$. This is given in Fig. 7 where m_A is the mass of the heavy MSSM CP-odd scalar that, at the order we are working at, is equal to $M_{H'}$.¹² Stop left-right mixing effects or extra D -term effects can be easily included along the lines of [48]. Corrections to g_{3h} can also arise at order $O(v^2/M_{H'}^2)$, but we already said that this coupling is difficult to measure as it requires double Higgs production.

6.2 The Minimal Composite Higgs Model (MCHM)

For models in which the Higgs is a pseudo Goldstone boson (PGB) arising from a new strong-sector at the TeV [49], similar to a pion in QCD, the Higgs couplings must depart from their SM value. This was studied in generality in [4]. The main effects are expected to arise in the Higgs coupling to Z/W and fermions. The minimal model is the MCHM [50], where the global symmetry-breaking pattern is $SO(5) \rightarrow SO(4)$ with an "order parameter" f , that give the following predictions [4]:

$$\begin{aligned} \frac{g_{VV}^h}{g_{VV}^{h, \text{SM}}} &= \sqrt{1 - \frac{v^2}{f^2}}, \\ \frac{g_{ff}^h}{g_{ff}^{h, \text{SM}}} &= \frac{1 - (1+n)v^2/f^2}{\sqrt{1 - v^2/f^2}}, \end{aligned} \quad (49)$$

where $n = 0, 1, 2, \dots$ depends on how fermions are implemented in the model. In particular, for the MCHM4 (MCHM5) we have $n = 0$ (1) [45]. From the minimization of the Higgs potential, we expect $f \gtrsim v$ [4], but constraints from the \hat{S} parameter give $v^2/f^2 \lesssim 0.1$ [49]. The Higgs coupling predictions of the MCHM are shown in Fig. 8 and compared with a fit of the ATLAS data. The fact that the experimental data does not favour smaller Higgs couplings than those of the SM, as predicted from Eq. (49), implies that we can derive an upper bound on $\xi \equiv v^2/f^2$, and consequently on the composite scale, $\Lambda \simeq g_* f$, where g_* is here the coupling among the resonances of the strong sector, expected to be in the range, $1 \ll g_* \lesssim 4\pi$. ATLAS [47] gives the observed (expected) 95% CL upper limit of $\xi < 0.12$ (0.29) for the MCHM4 and $\xi < 0.15$ (0.20) for the MCHM5 that start being as competitive as the ones coming from LEP [49].

¹²Mass splittings among the heavy Higgs-doublet components are $O(v^2/M_{H'}^2)$, and then their effects are of higher-order in our expansion.

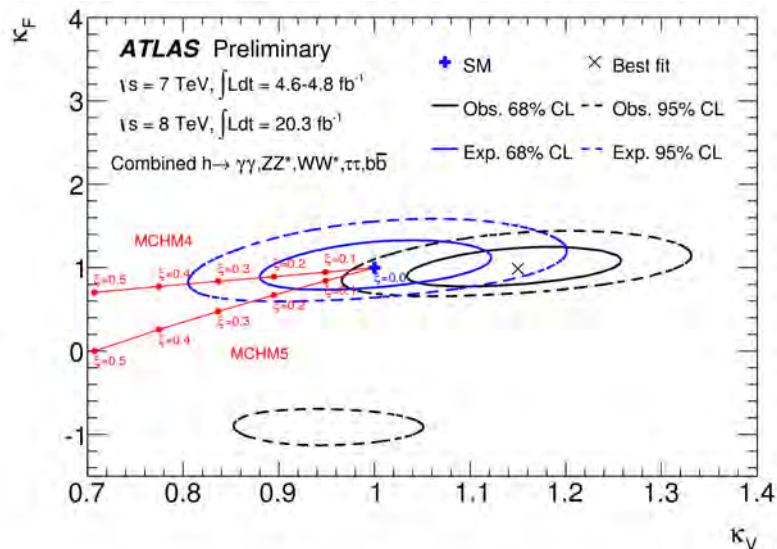


Fig. 8: Two-dimensional fit of the Higgs couplings $\kappa_V \equiv g_{VV}^h/g_{VV}^{h,SM}$ and $\kappa_F \equiv g_{ff}^h/g_{ff}^{h,SM}$ and predictions from the MCHM4 and MCHM5 as a function of $\xi \equiv v^2/f^2$ [47].

Contributions to $\kappa_{\gamma\gamma}$ and κ_{GG} are suppressed in the MCHM due to the PGB nature of the Higgs [4]. Nevertheless, this suppression is not present in $\kappa_{Z\gamma}$ that can receive significant contributions [51] that could be even larger than those of the SM, providing a strong motivation for searching for $h \rightarrow Z\gamma$.

7 Conclusions

With the Higgs discovery, the full SM has been experimentally established. Nevertheless, the presence of the Higgs, a zero-spin state, demands new physics at the TeV to make the SM a natural theory. The Higgs is the most sensitive SM particle to new physics, and for this reason an accurate measurement of its couplings provides an excellent way to indirectly discover new phenomena.

At the LHC (and in future colliders) we can have access to a large variety of Higgs couplings. We have argued that the most relevant Higgs couplings are the primary ones, given in Eq. (2) for CP-conservation. These couplings probe new directions in the parameter space of BSMs. We have showed the predictions for two of the most well-motivated BSMs, the MSSM and the MCHM. These analysis can be extended to other BSMs, such as the non-minimal MSSM (NMSSM), or other possibilities for composite Higgs, for example those in which the Higgs is lighter than the composite scale Λ not because of its PGB nature, as in the MCHM, but due to an "accidental" supersymmetry (SUSY Composite Higgs) or scale symmetry (Higgs as a dilaton) [49]. Supersymmetry can also allow for partly-composite Higgs where the TeV strong-sector could also break the electroweak symmetry (bosonic TC) [24]. A brief summary of the largest effects in the primary Higgs coupling arising from these scenarios is giving in Table 1. If in the future departures from the SM Higgs couplings are observed, the analysis of the pattern of these deviations will be extremely useful to discriminate between different BSM scenarios.

Acknowledgements

I am very thankful to Christophe Grojean, Eduard Masso and Francesco Riva for useful comments on the manuscript. This work has been partly supported by the grants FPA2011-25948, 2009SGR894 and the Catalan ICREA Academia Program.

	g_{ff}^h	g_{VV}^h	κ_{GG}	$\kappa_{\gamma\gamma}$	$\kappa_{Z\gamma}$	g_{3h}
MSSM	✓					✓
NMSSM	✓	✓	✓	✓	✓	✓
MCHM	✓	✓			✓	✓
SUSY Composite Higgs	✓	✓	✓	✓	✓	✓
Higgs as a Dilaton			✓	✓	✓	✓
Partly-Composite Higgs			✓	✓	✓	✓
Bosonic TC						✓

Table 1: Largest contributions to Higgs couplings (relative to the SM one) expected from different BSM scenarios.

References

- [1] G. Aad *et al.*, [ATLAS Collab.], Phys. Lett. B **716** (2012) 1; S. Chatrchyan *et al.*, [CMS Collab.], Phys. Lett. B **716** (2012) 30.
- [2] A. David for the CMS Collaboration, Talk at ICHEP14 (Valencia); see also CMS-PAS-HIG-14-009.
- [3] M. Kado for the ATLAS Collaboration, Talk at ICHEP14 (Valencia); see also ATLAS-CONF-2014-009, ATLAS-COM-CONF-2014-013.
- [4] G. F. Giudice, C. Grojean, A. Pomarol and R. Rattazzi, JHEP **0706** (2007) 045 [hep-ph/0703164].
- [5] S. Weinberg, Phys. Rev. D **19** (1979) 1277; L. Susskind, Phys. Rev. D **20** (1979) 2619; P. Sikivie, L. Susskind, M. B. Voloshin and V. I. Zakharov, Nucl. Phys. B **173** (1980) 189.
- [6] Work in preparation.
- [7] We follow the conventions of M. E. Peskin and D. V. Schroeder, “An Introduction to quantum field theory,” Reading, USA: Addison-Wesley (1995) 842 p.
- [8] R. Contino, M. Ghezzi, C. Grojean, M. Muhlleitner and M. Spira, JHEP **1307** (2013) 035 [arXiv:1303.3876 [hep-ph]].
- [9] J. Elias-Miro, J. R. Espinosa, E. Masso and A. Pomarol, JHEP **1311** (2013) 066 [arXiv:1308.1879 [hep-ph]].
- [10] J. F. Gunion, H. E. Haber, G. L. Kane and S. Dawson, Front. Phys. **80** (2000) 1.
- [11] B. Grzadkowski, M. Iskrzynski, M. Misiak and J. Rosiek, JHEP **1010** (2010) 085 [arXiv:1008.4884 [hep-ph]].
- [12] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. **65** (1990) 964.
- [13] R. S. Gupta, A. Pomarol and F. Riva, arXiv:1405.0181 [hep-ph].
- [14] E. Masso, arXiv:1406.6376 [hep-ph].
- [15] M. B. Einhorn and J. Wudka, Nucl. Phys. B **876** (2013) 556 [arXiv:1307.0478 [hep-ph]].
- [16] K. Hagiwara, S. Ishihara, R. Szalapski and D. Zeppenfeld, Phys. Lett. B **283** (1992) 353; Phys. Rev. D **48** (1993) 2182.
- [17] J. Elias-Miro, J. R. Espinosa, E. Masso and A. Pomarol, JHEP **1308** (2013) 033 [arXiv:1302.5661 [hep-ph]].
- [18] C. Grojean, E. E. Jenkins, A. V. Manohar and M. Trott, JHEP **1304** (2013) 016 [arXiv:1301.2588 [hep-ph]]; E. E. Jenkins, A. V. Manohar and M. Trott, JHEP **1310** (2013) 087 [arXiv:1308.2627 [hep-ph]]; JHEP **1401** (2014) 035 [arXiv:1310.4838 [hep-ph]]; J. Elias-Miro, C. Grojean, R. S. Gupta and D. Marzocca, JHEP **1405** (2014) 019 [arXiv:1312.2928 [hep-ph]].
- [19] A. Pomarol and F. Riva, JHEP **1401** (2014) 151 [arXiv:1308.2803 [hep-ph]].

- [20] S. Schael *et al.* [ALEPH and DELPHI and L3 and OPAL and SLD and LEP Electroweak Working Group and SLD Electroweak Group and SLD Heavy Flavour Group Collaborations], Phys. Rept. **427** (2006) 257 [hep-ex/0509008].
- [21] The LEP collaborations ALEPH, DELPHI, L3, OPAL, and the LEP TGC Working Group, LEPEWWG/TGC/2003-01.
- [22] Model-independent bounds were obtained for the first time in Z. Han and W. Skiba, Phys. Rev. D **71** (2005) 075009 [hep-ph/0412166], and, more recently, in [19]; J. Ellis, V. Sanz and T. You, arXiv:1410.7703 [hep-ph]; A. Falkowski and F. Riva, arXiv:1411.0669 [hep-ph]. See also M. Ciuchini, E. Franco, S. Mishima and L. Silvestrini, JHEP **1308** (2013) 106 [arXiv:1306.4644 [hep-ph]]; M. Ciuchini, E. Franco, S. Mishima, M. Pierini, L. Reina and L. Silvestrini, arXiv:1410.6940 [hep-ph], though the bounds derived there are not fully model-independent.
- [23] R. Barbieri *et al.*, Nucl. Phys. B **703** (2004) 127 [hep-ph/0405040].
- [24] For models with this property, see for example, M. Dine, A. Kagan and S. Samuel, Phys. Lett. B **243** (1990) 250, and more recently, A. Azatov, J. Galloway and M. A. Luty, Phys. Rev. Lett. **108** (2012) 041802 [arXiv:1106.3346 [hep-ph]]; Phys. Rev. D **85** (2012) 015018 [arXiv:1106.4815 [hep-ph]]; T. Gherghetta and A. Pomarol, JHEP **1112** (2011) 069 [arXiv:1107.4697 [hep-ph]].
- [25] A. Manohar and H. Georgi, Nucl. Phys. B **234** (1984) 189.
- [26] See <https://twiki.cern.ch/twiki/bin/view/LHCPhysics/LHCHXSWG>.
- [27] CMS Collaboration, arXiv:1307.5515 [hep-ex]; ATLAS Collaboration, ATLAS-CONF-2013-009.
- [28] A. Falkowski, F. Riva and A. Urbano, JHEP **1311** (2013) 111 [arXiv:1303.1812 [hep-ph]].
- [29] C. Englert, A. Freitas, M. M. Muhlleitner, T. Plehn, M. Rauch, M. Spira and K. Walz, J. Phys. G **41** (2014) 113001 [arXiv:1403.7191 [hep-ph]].
- [30] R. Contino, C. Grojean, M. Moretti, F. Piccinini and R. Rattazzi, JHEP **1005** (2010) 089 [arXiv:1002.1011 [hep-ph]].
- [31] CMS Collaboration, CMS-PAS-HIG-14-005 (2014).
- [32] CMS Collaboration, CMS-PAS-HIG-13-002.
- [33] C. Delaunay, G. Perez, H. de Sandes and W. Skiba, Phys. Rev. D **89** (2014) 035004 [arXiv:1308.4930 [hep-ph]]; F. Bishara, Y. Grossman, R. Harnik, D. J. Robinson, J. Shu and J. Zupan, JHEP **1404** (2014) 084 [arXiv:1312.2955 [hep-ph]].
- [34] D. McKeen, M. Pospelov and A. Ritz, Phys. Rev. D **86** (2012) 113004 [arXiv:1208.4597 [hep-ph]].
- [35] J. Ellis, V. Sanz and T. You, JHEP **1407** (2014) 036 [arXiv:1404.3667 [hep-ph]].
- [36] A. Biekötter, A. Knochel, M. Kraemer, D. Liu and F. Riva, arXiv:1406.7320 [hep-ph].
- [37] G. Aad *et al.* [ATLAS Collaboration], Phys. Rev. D **87** (2013) 11, 112001 [Erratum-ibid. D **88** (2013) 7, 079906] [arXiv:1210.2979 [hep-ex]].
- [38] See, for example, S. Chang, R. Dermisek, J. F. Gunion and N. Weiner, Ann. Rev. Nucl. Part. Sci. **58** (2008) 75 [arXiv:0801.4554 [hep-ph]].
- [39] F. Riva, C. Biggio and A. Pomarol, JHEP **1302** (2013) 081 [arXiv:1211.4526 [hep-ph]].
- [40] M. Gonzalez-Alonso and G. Isidori, Phys. Lett. B **733** (2014) 359 [arXiv:1403.2648 [hep-ph]]; A. Falkowski and R. Vega-Morales, arXiv:1405.1095 [hep-ph].
- [41] See, for example, A. Djouadi, Phys. Rept. **459** (2008) 1 [hep-ph/0503173].
- [42] R. Contino, L. Da Rold and A. Pomarol, Phys. Rev. D **75** (2007) 055014 [hep-ph/0612048].
- [43] See, for example, L. J. Hall, D. Pinner and J. T. Ruderman, JHEP **1204** (2012) 131 [arXiv:1112.2703 [hep-ph]].
- [44] O. Matsedonskyi, G. Panico and A. Wulzer, JHEP **1301** (2013) 164 [arXiv:1204.6333 [hep-ph]]. M. Redi and A. Tesi, JHEP **1210** (2012) 166 [arXiv:1205.0232 [hep-ph]]; D. Marzocca, M. Serone and J. Shu, JHEP **1208** (2012) 013 [arXiv:1205.0770 [hep-ph]].

- [45] A. Pomarol and F. Riva, *JHEP* **1208** (2012) 135 [arXiv:1205.6434 [hep-ph]].
- [46] See, for example, A. Arvanitaki and G. Villadoro, *JHEP* **1202** (2012) 144 [arXiv:1112.4835 [hep-ph]]; A. Azatov, S. Chang, N. Craig and J. Galloway, *Phys. Rev. D* **86** (2012) 075033 [arXiv:1206.1058 [hep-ph]].
- [47] ATLAS Collaboration, ATLAS-CONF-2014-010.
- [48] R. S. Gupta, M. Montull and F. Riva, *JHEP* **1304** (2013) 132 [arXiv:1212.5240 [hep-ph]].
- [49] For a review, see for example, R. Contino, arXiv:1005.4269 [hep-ph]; B. Bellazzini, C. Csáki and J. Serra, *Eur. Phys. J. C* **74** (2014) 2766 [arXiv:1401.2457 [hep-ph]].
- [50] K. Agashe, R. Contino and A. Pomarol, *Nucl. Phys. B* **719** (2005) 165 [hep-ph/0412089].
- [51] A. Azatov, R. Contino, A. Di Iura and J. Galloway, *Phys. Rev. D* **88** (2013) 7, 075019 [arXiv:1308.2676 [hep-ph]].