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Effective field theory approach to LHC Higgs data

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Abstract. The effective field theory approach to LHC Higgs data is reviewed.

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1. Introduction

The Standard Model (SM) of particle physics was proposed back in the 1960s as a theory of quarks and leptons interacting via strong, weak, and electromagnetic forces [1]. It is built on the following principles:

(1) The basic framework is that of a relativistic quantum field theory, with interactions between particles described by a local Lagrangian.
(2) The Lagrangian is invariant under the linearly realized local $SU(3) \times SU(2) \times U(1)$ symmetry.
(3) The vacuum state of the theory preserves only $SU(3) \times U(1)$ local symmetry, as a result of the Brout–Englert–Higgs mechanism [2–4]. The spontaneous breaking of the $SU(2) \times U(1)$ symmetry down to $U(1)$ arises due to the vacuum expectation value (VEV) of a scalar field transforming as $(1, 2)_{1/2}$ under the local symmetry.
(4) Interactions are renormalizable, which means that only interactions up to the canonical mass dimension-four are allowed in the Lagrangian.

Given the experimentally observed matter content (three families of quarks and leptons), these rules completely specify the theory up to 19 free parameters. The local symmetry implies the presence of spin-1 vector bosons which mediate the strong and electroweak forces. The breaking pattern of the local symmetry ensures that the carriers of the strong and electromagnetic forces are massless, whereas the carriers of the weak force are massive. Finally, the particular realization of the Brout–Englert–Higgs mechanism in the SM leads to the emergence of exactly one spin-0 scalar boson – the famous Higgs boson [5–7].

The SM passed an incredible number of experimental tests. It correctly describes the rates and differential distributions of particles produced in high-energy collisions; a robust deviation from the SM predictions has never been observed. It very accurately predicts many properties of elementary particles, such as the magnetic and electric dipole moments, as well as certain properties of simple enough composite particles, such as atomic energy levels. The discovery of a 125 GeV boson at the Large Hadron Collider (LHC) [8,9] nails down the last propagating degree of freedom predicted by the SM. Measurements of its production and decay rates vindicate the simplest realization of the Brout–Englert–Higgs mechanism, in which a VEV of a single $SU(2)$ doublet field spontaneously breaks the electroweak symmetry. Last but not the least, the SM is a consistent quantum theory (as long as the gravitational interactions can be neglected). In particular, for the measured value of the Higgs boson mass the vacuum of the theory is metastable, with a lifetime which is many orders of magnitude longer than the age of the Universe. Therefore, the validity range of the SM can be extended all the way up to the Planck scale (at which point the gravitational interactions become strong and can no longer be neglected) without encountering any theoretical inconsistency.

Yet we know that the SM is not the ultimate theory. It cannot account for dark matter, neutrino masses, matter/antimatter asymmetry, and cosmic inflation, all of which are experimental facts. In addition, some theoretical or aesthetic arguments (the strong CP problem, flavour hierarchies, unification, the naturalness problem) suggest that the SM should be extended. This justifies the ongoing searches for new physics, that is particles or interactions not predicted by the SM.

In spite of good arguments for the existence of new physics, a growing body of evidence suggests that, at least up to energies of a few hundred GeV, the fundamental degrees of freedom are those of the SM. Given the absence of any direct or indirect collider signal of new physics, it is reasonable to assume that new particles from beyond the SM are much heavier than the SM particles. If that is correct, physics at the weak scale can be adequately described using effective field theory (EFT) methods.

In the EFT framework adopted here, the assumptions $1, \ldots, 3$ continue to be valid [9a]. Thus, much as in the SM, the Lagrangian is constructed from gauge-invariant operators involving the SM fermion, gauge, and Higgs fields. The difference is that assumption 4 is dropped and interactions with arbitrary large mass dimension D are allowed. These interactions can be organized in a systematic expansion in D. The leading-order term in this expansion is the SM Lagrangian with operators up to $D = 4$. All possible effects of heavy new physics are encoded in operators with $D > 4$, which are suppressed in the Lagrangian by appropriate powers of the mass scale Λ. Since all $D = 5$ operators violate lepton number and are thus stringently constrained by the experiment, the leading corrections to the Higgs observables are expected from $D = 6$ operators suppressed by Λ^2 [14,14a]. It is assumed that the operators with $D > 6$ can be ignored, which is always true for $v \ll \Lambda$.

This review discusses the interpretation of the LHC data on the Higgs boson production and decay in the framework of an EFT beyond the SM. For practical reasons, three more assumptions about higher-dimensional operators are adopted:

(5) The baryon and lepton numbers are conserved.
(6) The coefficients of operators involving fermions are flavour conserving and universal, except for Yukawa-type operators, which are aligned with the corresponding SM Yukawa matrices [14a].
(7) The corrections from $D = 6$ operators to the Higgs signal strength are subleading compared to the SM contribution.

Other than that, the discussion will be model-independent.

In the following section the SM Lagrangian is reviewed, in order to prepare the ground and fix the notation. The part of the $D = 6$ effective Lagrangian relevant for Higgs studies is discussed in §3. The dependence of the Higgs signal strength measured at the LHC on the effective Lagrangian parameters is summarized in §4. The experimental results and the current model-independent constraints on the $D = 6$ parameters are discussed in §5. The bibliography contains a number of references where an EFT-inspired approach to physics of the 125 GeV Higgs at the LHC is exercised.

2. Standard Model Lagrangian

Here, the SM Lagrangian is summarized and notations are defined.

The SM Lagrangian is invariant under the global Poincaré symmetry (Lorentz symmetry + translations) and a local symmetry with the gauge group $G_{\rm SM} = SU(3)_C \times SU(2)_L \times U(1)_Y$. The fields building the SM Lagrangian fill representations of these symmetries. The field content of the SM is the following:

(a) Vector fields G^a_μ, W^i_μ, B_μ, where $i = 1, \ldots, 3$ and $a = 1, \ldots, 8$. They transform as four-vectors under the Lorentz symmetry and are the gauge fields of the $G_{\rm SM}$ group.
(b) Three generations of fermionic fields $q = (u, V_{\rm CKM} d)$, u^c, d^c, $\ell = (\nu, e)$, e^c. They transform as 2-component spinors under the Lorentz symmetry [18a]. The transformation properties under $G_{\rm SM}$ are listed in table 1.
(c) Scalar field $H = (H^+, H^0)$ transforming as $(1, 2)_{1/2}$ under $G_{\rm SM}$. $\tilde{H}_i = \epsilon_{ij} H^*_j$ that transforms as $(1, 2)_{-1/2}$ is also defined.

The SM Lagrangian can be split as

$$\mathcal{L}^{\rm SM} = \mathcal{L}^{\rm SM}_V + \mathcal{L}^{\rm SM}_F + \mathcal{L}^{\rm SM}_H + \mathcal{L}^{\rm SM}_Y. \tag{2.1}$$

Table 1. Representation of the SM scalar and fermion fields under the SM gauge group.

	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
$q = \begin{pmatrix} u \\ d \end{pmatrix}$	**3**	**2**	1/6
u^c	$\bar{\mathbf{3}}$	**1**	−2/3
d^c	$\bar{\mathbf{3}}$	**1**	1/3
$\ell = \begin{pmatrix} \nu \\ e \end{pmatrix}$	**1**	**2**	−1/2
e^c	**1**	**1**	1
H	**1**	**2**	1/2

The first term above contains gauge-invariant kinetic terms for the vector fields [19a]:

$$\mathcal{L}_V^{\rm SM} = -\frac{1}{4g_s^2} G^a_{\mu\nu} G^a_{\mu\nu} - \frac{1}{4g_L^2} W^i_{\mu\nu} W^i_{\mu\nu} - \frac{1}{4g_Y^2} B_{\mu\nu} B_{\mu\nu}, \tag{2.2}$$

where g_s, g_L, g_Y are gauge couplings of $SU(3)_C \times SU(2)_L \times U(1)_Y$, here defined as the normalization of the appropriate gauge kinetic term. The electromagnetic coupling $e = g_L g_Y/\sqrt{g_L^2 + g_Y^2}$ and the weak mixing angle $s_\theta = g_Y/\sqrt{g_L^2 + g_Y^2}$ are also defined. The field strength tensors are given by

$$\begin{aligned} B_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu, \\ W^i_{\mu\nu} &= \partial_\mu W^i_\nu - \partial_\nu W^i_\mu + \epsilon^{ijk} W^j_\mu W^k_\nu, \\ G^a_{\mu\nu} &= \partial_\mu G^a_\nu - \partial_\nu G^a_\mu + f^{abc} G^b_\mu G^c_\nu, \end{aligned} \tag{2.3}$$

where ϵ^{ijk} and f^{abc} are the totally antisymmetric structure tensors of $SU(2)$ and $SU(3)$.

The second term in eq. (2.1) contains covariant kinetic terms of the fermion fields:

$$\begin{aligned} \mathcal{L}_F^{\rm SM} &= i\bar{q}\bar{\sigma}_\mu D_\mu q + iu^c \sigma_\mu D_\mu \bar{u}^c + id^c \sigma_\mu D_\mu \bar{d}^c \\ &\quad + i\bar{\ell}\bar{\sigma}_\mu D_\mu \ell + ie^c \sigma_\mu D_\mu \bar{e}^c. \end{aligned} \tag{2.4}$$

Each fermion field is a 3-component vector in the generation space. It is assumed that all the rotations needed to put fermions in the mass eigenstate basis have already been made; in the SM the only residue of these rotations is the CKM matrix appearing in the definition of the quark doublet components. The covariant derivatives are defined as

$$D_\mu f = (\partial_\mu - iG^a_\mu T^a_f - iW^i_\mu T^i_f - iY_f B_\mu) f. \tag{2.5}$$

Here $T^a_f = (\lambda^a, -\lambda^a, 0)$ for f in the triplet/antitriplet/singlet representation of $SU(3)$, where λ^a are Gell-Mann matrices; $T^i_f = (\sigma^i/2, 0)$ for f in the doublet/singlet representation of $SU(2)$; Y_f is the $U(1)$ hypercharge. The electric charge is given by $Q_f = T^3_f + Y_f$.

The third term in eq. (2.1) contains Yukawa interactions between the Higgs field and the fermions:

$$\mathcal{L}_Y^{\rm SM} = -\tilde{H}^\dagger u^c y_u q - H^\dagger d^c y_d q - H^\dagger e^c y_e \ell + \text{h.c.}, \tag{2.6}$$

where y_f are 3×3 diagonal matrices.

The last term in eq. (2.1) are the Higgs kinetic and potential terms:

$$\mathcal{L}_H^{\rm SM} = D_\mu H^\dagger D_\mu H + \mu_H^2 H^\dagger H - \lambda (H^\dagger H)^2, \tag{2.7}$$

where the covariant derivative acting on the Higgs field is

$$D_\mu H = \left(\partial_\mu - \frac{i}{2} W^i_\mu \sigma^i - \frac{i}{2} B_\mu\right) H. \tag{2.8}$$

Because of the negative mass squared term μ_H^2 in the Higgs potential the Higgs field gets a VEV,

$$\langle H \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}, \quad \mu_H^2 = \lambda v^2. \tag{2.9}$$

This generates mass terms for W^i_μ and B_μ and a field rotation is needed to diagonalize the mass matrix. The mass eigenstates are defined to the electroweak vector fields by

$$\begin{aligned} W^1_\mu &= \frac{g_L}{\sqrt{2}} \left(W^+_\mu + W^-_\mu\right), \\ W^3_\mu &= \frac{g_L}{\sqrt{g_L^2 + g_Y^2}} \left(g_L Z_\mu + g_Y A_\mu\right), \\ W^2_\mu &= \frac{ig_L}{\sqrt{2}} \left(W^+_\mu - W^-_\mu\right), \\ B_\mu &= \frac{g_Y}{\sqrt{g_L^2 + g_Y^2}} \left(-g_Y Z_\mu + g_L A_\mu\right). \end{aligned} \tag{2.10}$$

The mass eigenstates are defined such that their quadratic terms are canonically normalized and their mass terms are diagonal:

$$\begin{aligned} \mathcal{L}_{V,\rm kin}^{\rm SM} &= -\frac{1}{2} W^+_{\mu\nu} W^-_{\mu\nu} - \frac{1}{4} Z_{\mu\nu} Z_{\mu\nu} - \frac{1}{4} A_{\mu\nu} A_{\mu\nu} \\ &\quad + m_W^2 W^+_\mu W^-_\mu + \frac{m_Z^2}{2} Z_\mu Z_\mu, \end{aligned} \tag{2.11}$$

where the W and Z boson masses are

$$m_W = \frac{g_L v}{2}, \quad m_Z = \frac{\sqrt{g_L^2 + g_Y^2}\, v}{2}. \tag{2.12}$$

The SM fermions (except for the neutrinos) also acquire masses after electroweak symmetry breaking via the Yukawa interactions in eq. (2.6). Here, a basis in the fermion flavour space is chosen where the Yukawa interactions are diagonal, in which case the fermion masses are given by

$$m_{f_i} = \frac{v}{\sqrt{2}} [y_f]_{ii}. \tag{2.13}$$

Interactions of the gauge boson mass eigenstates with fermions are given by

$$\begin{aligned}\mathcal{L}_{vff}^{\rm SM} &= eA_\mu \sum_{f\in u,d,e} Q_f(\bar{f}\bar{\sigma}_\mu f + f^c\sigma_\mu \bar{f}^c) \\ &+ g_s G_\mu^a \sum_{f\in u,d} (\bar{f}\bar{\sigma}_\mu T^a f + f^c\sigma_\mu T^a \bar{f}^c) \\ &+ \frac{g_L}{\sqrt{2}}\left(W_\mu^+ \bar{u}\bar{\sigma}_\mu V_{\rm CKM} d + W_\mu^+ \bar{\nu}\bar{\sigma}_\mu e + {\rm h.c.}\right) \\ &+ \sqrt{g_L^2+g_Y^2} Z_\mu \left[\sum_{f\in u,d,e,\nu} \bar{f}\bar{\sigma}_\mu (T_f^3 - s_\theta^2 Q_f) f \right. \\ &\left. + \sum_{f\in u,d,e} f^c\sigma_\mu(-s_\theta^2 Q_f)\bar{f}^c\right]. \end{aligned} \tag{2.14}$$

Finally, let us move to the Higgs sector. After electroweak symmetry breaking, the Higgs doublet field can be conveniently written as

$$H = \frac{1}{\sqrt{2}}\begin{pmatrix} \sqrt{2}G^+ \\ v + h + iG^0 \end{pmatrix}. \tag{2.15}$$

The fields G^0 and G^+ do not correspond to new physical degrees of freedom (they kinetically mix with the massive gauge bosons and can be gauged away). From now on, I will work in the unitary gauge and set $G^\pm = 0 = G^0$. The star of this review – the scalar field h – is called the Higgs boson. Its mass can be expressed by the parameters of the Higgs potential as

$$m_h^2 = 2\mu_H^2 = 2\lambda v^2. \tag{2.16}$$

The interactions in the SM Lagrangian involving a single Higgs boson are the following

$$\begin{aligned}\mathcal{L}_h^{\rm SM} &= \frac{h}{v}\left[\frac{g_L^2 v^2}{2} W_\mu^+ W_\mu^- + \frac{(g_L^2+g_Y^2)v^2}{4} Z_\mu Z_\mu\right] \\ &- \frac{h}{v}\sum_f m_f\left(ff^c + {\rm h.c.}\right). \end{aligned} \tag{2.17}$$

Roughly speaking, the Higgs boson couples to mass, in the sense that it couples to pairs of SM particles with the strength proportional to their masses (for fermions) or masses squared (for bosons). As all the masses have been measured by experiment, the strength of Higgs boson interactions in the SM is precisely predicted and contains no free parameters.

This section is concluded with a summary of the SM parameters used in this review. For the Higgs boson mass, $m_h = 125.09$ GeV is taken, which is the central value of the recent ATLAS and CMS combination of mass, measurements [20]. The gauge boson masses are $m_W = 80.385$ GeV [21] and $m_Z = 91.1875$ GeV [22]. The Higgs VEV is calculated from the muon lifetime (equivalently, from the Fermi constant $G_{\rm F} = 1/\sqrt{2}v^2 = 1.16637 \times 10^{-5}$ GeV^{-2} [23]), corresponding to $v = 246.221$ GeV. The electroweak couplings at the Z boson mass scale are extracted from m_Z and the electromagnetic structure constant $\alpha(m_Z) = 7.755 \times 10^{-3}$ [24], and the strong coupling from $\alpha_s(m_Z) = 1.172 \times 10^{-3}$ [23]. To evaluate corrections to the Higgs observables, the couplings up to the scale m_h are used: $g_s = 1.187$, $g_L = 0.643$, $g_Y = 0.358$. The light fermion masses are also evaluated at the scale m_h: the relevant ones are $m_b = 2.76$ GeV, $m_\tau = 1.78$ GeV, and $m_c = 0.62$ GeV. For the top mass, $m_t = 173.2$ GeV is taken.

3. Dimension-six Lagrangian

The effective Lagrangian of the form,

$$\mathcal{L}_{\rm eff} = \mathcal{L}^{\rm SM} + \mathcal{L}^{D=6}, \quad \mathcal{L}^{D=6} = \frac{1}{v^2}\sum_\alpha c_\alpha O_\alpha, \tag{3.1}$$

is considered where $\mathcal{L}^{\rm SM}$ is the SM Lagrangian discussed in §2 and O_α is a complete basis of $SU(3) \times SU(2) \times U(1)$ invariant $D = 6$ operators constructed out of the SM fields. In general, such a basis contains 2499 independent operators after imposing baryon and lepton number conservation [25]. One of the assumptions in this review is that coefficients of $D = 6$ operators are flavour-universal, which brings the number of independent parameters down to 76. Furthermore, only nine combinations of these operators will be relevant for a completely general description of the Higgs signal strength measurements considered later in this review.

One can choose a complete, non-redundant basis of operators in many distinct (though ultimately equivalent) ways. Here I work with the so-called *Higgs basis* introduced in ref. [26] and inspired by refs [27,28,28a]. The basis is spanned by particular combinations of $D = 6$ operators. Each of these combinations maps to an interaction term of the SM mass-eigenstates in the tree-level effective Lagrangian. The coefficients multiplying these combinations in the Lagrangian are called the independent couplings. The single Higgs couplings to pairs of gauge bosons and fermions are chosen among the independent couplings. The advantage of this basis is that the independent couplings are related in a simple way to observables in Higgs physics.

Most often, an $SU(3) \times SU(2) \times U(1)$ invariant operator gives rise to more than one interaction term of mass eigenstates. This leads to relations between various couplings in the effective Lagrangian. Therefore, several of these couplings are not free but can be expressed in terms of the independent couplings; they are called the dependent couplings. For example, at the level of the $D = 6$ Lagrangian, the W boson couplings to fermions are dependent couplings, as they can be expressed in terms of the Z boson couplings to fermions. Similarly, all the Higgs boson couplings to the W boson are dependent couplings, as they can be expressed via the Higgs boson couplings to Z and γ (see eq. (3.4)). Of course, the choice of which couplings are chosen as independent and which are dependent is subjective and dictated by convenience.

Now, the part of $D = 6$ Lagrangian in the Higgs basis that is relevant for LHC Higgs observables is reviewed; see ref. [26] for the full set of independent couplings and the algorithm to construct the complete $D = 6$ Lagrangian. In this formalism, by construction, all kinetic terms are canonically normalized, there is no kinetic mixing between the Z boson and the photon, and there is no correction to the Z boson mass term. While, in general, $D = 6$ operators do generate mixing and mass corrections, the canonical form can always be recovered by using equations of motion, integration by parts, and redefinition of fields and couplings. Thus, the kinetic and mass terms for the electroweak gauge bosons are those in eq. (2.11), except for the correction to the W boson mass term: $\Delta\mathcal{L}^{D=6}_{\text{kinetic}} = 2\delta m(g^2v^2/4)W^+_\mu W^-_\mu$. The independent coupling δm is a free parameter from the EFT point of view. However, it is very well constrained by experiment: $\delta m = (2.6 \pm 1.9) \cdot 10^{-4}$ [32]. Given the precision of LHC data, effects proportional to δm are currently not relevant for Higgs searches and will be ignored.

Now, let us move to interactions of a single Higgs boson with pairs of SM gauge bosons and fermions. The SM interactions of this type are given in eq. (2.17) and they contain no free parameters. Dimension-six operators lead to shifts of the couplings in eq. (2.17), as well as to the appearance of 2-derivative Higgs couplings to gauge bosons. In the Higgs basis, these effects are parametrized by the following independent couplings:

$$\delta c_z,\ c_{zz},\ c_{z\Box},\ c_{\gamma\gamma},\ c_{z\gamma},\ c_{gg},\ \tilde{c}_{gg},\ \tilde{c}_{zz},\ \tilde{c}_{\gamma\gamma},\ \tilde{c}_{z\gamma},$$
$$\delta y^u,\ \delta y^d,\ \delta y^e,\ \sin\phi^u,\ \sin\phi^d,\ \sin\phi^\ell. \quad (3.2)$$

The couplings in the first line are defined via the Higgs boson couplings to gauge bosons:

$$\begin{aligned}\Delta\mathcal{L}^{D=6}_{hvv} = \frac{h}{v}\Bigg[&2\delta c_w m_W^2 W^+_\mu W^-_\mu + \delta c_z m_Z^2 Z_\mu Z_\mu \\ &+ c_{ww}\frac{g_L^2}{2}W^+_{\mu\nu}W^-_{\mu\nu} + \tilde{c}_{ww}\frac{g_L^2}{2}W^+_{\mu\nu},\tilde{W}^-_{\mu\nu} \\ &+ c_{w\Box}g_L^2\left(W^-_\mu\partial_\nu W^+_{\mu\nu} + \text{h.c.}\right) \\ &+ c_{gg}\frac{g_s^2}{4}G^a_{\mu\nu}G^a_{\mu\nu} + c_{\gamma\gamma}\frac{e^2}{4}A_{\mu\nu}A_{\mu\nu} \\ &+ c_{z\gamma}\frac{eg_L}{2c_\theta}Z_{\mu\nu}A_{\mu\nu} + c_{zz}\frac{g_L^2}{4c_\theta^2}Z_{\mu\nu}Z_{\mu\nu} \\ &+ c_{z\Box}g_L^2 Z_\mu\partial_\nu Z_{\mu\nu} + c_{\gamma\Box}g_L g_Y Z_\mu\partial_\nu A_{\mu\nu} \\ &+ \tilde{c}_{gg}\frac{g_s^2}{4}G^a_{\mu\nu}\tilde{G}^a_{\mu\nu} + \tilde{c}_{\gamma\gamma}\frac{e^2}{4}A_{\mu\nu}\tilde{A}_{\mu\nu} \\ &+ \tilde{c}_{z\gamma}\frac{eg_L}{2c_\theta}Z_{\mu\nu}\tilde{A}_{\mu\nu} + \tilde{c}_{zz}\frac{g_L^2}{4c_\theta^2}Z_{\mu\nu}\tilde{Z}_{\mu\nu}\Bigg], \end{aligned} \quad (3.3)$$

where the dependent couplings δc_w, c_{ww}, $\tilde{c}_{ww}$, $c_{w\Box}$, and $c_{\gamma\Box}$ can be expressed by the independent couplings as

$$\begin{aligned}\delta c_w &= \delta c_z + 4\delta m, \\ c_{ww} &= c_{zz} + 2s_\theta^2 c_{z\gamma} + s_\theta^4 c_{\gamma\gamma}, \\ \tilde{c}_{ww} &= \tilde{c}_{zz} + 2s_\theta^2\tilde{c}_{z\gamma} + s_\theta^4\tilde{c}_{\gamma\gamma}, \\ c_{w\Box} &= \frac{1}{g_L^2 - g_Y^2}\Big[g_L^2 c_{z\Box} + g_Y^2 c_{zz} - e^2 s_\theta^2 c_{\gamma\gamma} \\ &\quad - (g_L^2 - g_Y^2)s_\theta^2 c_{z\gamma}\Big], \\ c_{\gamma\Box} &= \frac{1}{g_L^2 - g_Y^2}\Big[2g_L^2 c_{z\Box} + (g_L^2 + g_Y^2)c_{zz} \\ &\quad - e^2 c_{\gamma\gamma} - (g_L^2 - g_Y^2)c_{z\gamma}\Big]. \end{aligned} \quad (3.4)$$

The coupling in the second line of eq. (3.2) are defined via the Higgs boson couplings to fermions:

$$\Delta\mathcal{L}^{D=6}_{hff} = -\frac{h}{v}\sum_{f\in u,d,e}\delta y_f\,\mathrm{e}^{i\phi_f}\,m_f f^c f + \text{h.c.} \quad (3.5)$$

Following my assumption of flavour universal coefficients of dimension-6 operators, each δy_f and ϕ_f is a real number. Moreover, the couplings in eq. (3.5) are diagonal in the generation space, and therefore flavour-violating Higgs decays are absent (see refs [33,34] for a discussion of such decays in the EFT language).

The complete Higgs interaction Lagrangian relevant for this review is given by $\mathcal{L}_h^{\text{SM}} + \mathcal{L}_{vff}^{\text{SM}} + \Delta\mathcal{L}^{D=6}_{hvv} + \Delta\mathcal{L}^{D=6}_{hff}$ and is parametrized by the independent couplings in eq. (3.2). The effect of these couplings on the LHC Higgs observables will be discussed in the following sections. But before that, a comment is in order on other effects of $D = 6$ operators that could, *a priori*, be relevant. First, in the Higgs basis there are corrections to the Z and W boson interactions in eq. (2.14), parametrized by vertex corrections δg. These would feed indirectly into Higgs observables, such as, for example, the vector boson fusion (VBF) production cross-section or the $h \to VV^* \to 4$ fermion decays. However, there are model-independent constraints on these vertex corrections [32] which ensure that their effects on Higgs observables are too small to be currently observable. For this reason, the vertex corrections are ignored in this review. Next, $D = 6$ operators may induce two classes of Higgs boson interactions that could affect $h \to VV^* \to 4$ fermion decays. One class is the $hVff$ vertex-like contact interactions:

$$\begin{aligned}\Delta\mathcal{L}_{hVff}^{D=6} &= \sqrt{2}g_L\frac{h}{v}W_\mu^+\Big(\delta g_L^{W\ell}\bar{\nu}\bar{\sigma}_\mu e + \delta g_L^{Wq}\bar{u}\bar{\sigma}_\mu d \\ &\quad + \delta g_R^{Wq}u^c\sigma_\mu\bar{d}^c + \text{h.c.}\Big) \\ &\quad + 2\frac{h}{v}\sqrt{g_L^2+g_Y^2}Z_\mu\left[\sum_{f=u,d,e,\nu}\delta g_L^{Zf}\bar{f}\bar{\sigma}_\mu f\right. \\ &\quad \left. + \sum_{f=u,d,e}\delta g_R^{Zf}f^c\sigma_\mu\bar{f}^c\right]. \end{aligned} \tag{3.6}$$

In the Higgs basis, the parameters δg above are equal to the corresponding vertex corrections to the SM couplings in eq. (2.14). Given the constraints on the δgs in ref. [32], the LHC Higgs studies cannot be currently sensitive to the vertex-like Higgs interactions, and therefore they are neglected in this analysis. The other class is the dipole-like contact interactions:

$$\begin{aligned}\Delta\mathcal{L}_{hdVff}^{D=6} &= -\frac{h}{4v^2}\left[g_s\sum_{f\in u,d}d_{Gf}f^c\bar{\sigma}_{\mu\nu}T^a f G_{\mu\nu}^a\right. \\ &\quad + e\sum_{f\in u,d,e}d_{Af}f^c\bar{\sigma}_{\mu\nu}fA_{\mu\nu} \\ &\quad + \sqrt{g_L^2+g_Y^2}\sum_{f\in u,d,e}d_{Zf}f^c\bar{\sigma}_{\mu\nu}fZ_{\mu\nu} \\ &\quad \left. + \sqrt{2}g_L d_{Wq}d^c\bar{\sigma}_{\mu\nu}uW_{\mu\nu}^- + \text{h.c.}\right] \\ &\quad - \frac{h}{4v^2}\left[\sum_{f\in u,d}\tilde{d}_{Gf}f^c\bar{\sigma}_{\mu\nu}T^a f\widetilde{G}_{\mu\nu}^a\right. \\ &\quad + e\sum_{f\in u,d,e}\tilde{d}_{Af}f^c\bar{\sigma}_{\mu\nu}f\widetilde{A}_{\mu\nu} \\ &\quad + \sqrt{g_L^2+g_Y^2}\sum_{f\in u,d,e}\tilde{d}_{Zf}f^c\bar{\sigma}_{\mu\nu}f\widetilde{Z}_{\mu\nu} \\ &\quad \left. + \sqrt{2}g_L\tilde{d}_{Wq}d^c\bar{\sigma}_{\mu\nu}u\widetilde{W}_{\mu\nu}^- + \text{h.c.}\right]. \end{aligned} \tag{3.7}$$

For Higgs decays into four light fermions, the dipole-like contributions do not interfere with the SM amplitudes due to different helicity structures. Therefore, corrections to the decay width enter quadratically in d_{Vf}, and should be neglected. Furthermore, as a consequence of the linearly realized electroweak symmetry in the $D = 6$ Lagrangian, the parameters d_{Vf} are proportional to the respective dipole moments which are stringently constrained by experiment, especially for light fermions. So, it is safe to neglect the dipole-like Higgs interactions for the sake of LHC analyses.

Finally, $D = 6$ operators produce several more interactions involving the single Higgs boson field, for example Higgs couplings to three gauge bosons. Observable effects of these couplings are extremely suppressed, and therefore they are not listed here. Moreover, new interactions involving two (or three) Higgs boson fields appear in the Lagrangian, and they are relevant for an EFT description of double Higgs production [35–41]. This review is focussed on single Higgs production, and therefore multi-Higgs couplings are not listed; see ref. [26] for the relevant expressions in the Higgs basis.

Let us close this section with a brief discussion of the validity range of this approach. Formally, EFT is an expansion in powers of the scale Λ suppressing higher-dimensional operators. As the independent couplings in eq. (3.2) arise from $D = 6$ operators, they are formally of order v^2/Λ^2. The rule of thumb is that the EFT approach to Higgs physics is valid if $\Lambda \gtrsim v$, which translates to $|c_i| \lesssim 1$ and $\delta y_f \lesssim v/m_f$ for the independent couplings. However, a detailed analysis of this issue is much more tricky and depends on the kinematic region probed by a given observable. For example, for observables probing the high $\sqrt{s}$ or high p_T tail of differential distributions the validity range will be different from inclusive observables (see ref. [42] for a more in-depth discussion of these issues). This review is restricted to the Higgs signal strength observables in various production modes, which are typically dominated by $\sqrt{s} \sim m_h$. Moreover, the question of the validity range is not discussed here because it is assumed from the onset that higher-dimensional operators provide small corrections on top of SM contributions. Consequently, only corrections to the observables are taken into account that are linear in the parameters in eq. (3.2), which corresponds to retaining only $\mathcal{O}(\Lambda^{-2})$ effects in the EFT expansion [42a]. Incidentally, the LHC so far confirms that the SM is a decent first approximation of the Higgs sector, and deviations due to new physics are small.

4. Observables

Consider the Higgs boson produced at the LHC via the process X, and subsequently decaying to the final state Y. It is possible, to an extent, to isolate experimentally different Higgs boson production modes and decay channels. The LHC Collaborations typically quote the Higgs signal strength relative to the SM one in a given channel, here denoted as $\mu_{X;Y}$. Thanks to the narrow width of the Higgs boson, the production and decay can be separated [42b]:

$$\mu_{X;Y} = \frac{\sigma(pp \to X)}{\sigma(pp \to X)_{\rm SM}} \frac{\Gamma(h \to Y)}{\Gamma(h \to Y)_{\rm SM}} \frac{\Gamma(h \to {\rm all})_{\rm SM}}{\Gamma(h \to {\rm all})}. \tag{4.1}$$

Now, how the Higgs production and decays depend on the parameters in the effective Lagrangian are summarized. These formulas allow one to derive experimental constraints on the EFT parameters. This kind of approach to LHC Higgs data was pioneered in refs [48,49] and perfected in refs [50–87]. As discussed at the end of §3, only linear corrections in the independent couplings are kept, while quadratic corrections are ignored. For this reason only CP-even couplings appear in these formulas (the CP-odd ones enter inclusive observables only at the quadratic level). Moreover, only $D = 6$ corrections are included at the *tree level* and new physics effects suppressed by a loop factor are neglected [88a]. The exception is the gluon fusion production process which is computed at the next-to-leading order in the $D = 6$ parameters. Unless noted otherwise, the inclusive production and decay rates are given here.

Production

For the relevant partonic processes of Higgs production at the LHC, the cross-section relative to the SM one depends on the effective theory parameters as follows:

(1) Gluon fusion (ggh), $gg \to h$:

$$\frac{\sigma_{ggh}}{\sigma_{ggh}^{\rm SM}} \simeq \left|1 + \frac{\hat{c}_{gg}}{c_{gg}^{\rm SM}}\right|^2, \tag{4.2}$$

where

$$\hat{c}_{gg} \simeq c_{gg} + \frac{1}{12\pi^2}\left[\delta y_u A_f\left(\frac{m_h^2}{4m_t^2}\right) + \delta y_d A_f\left(\frac{m_h^2}{4m_b^2}\right)\right],$$

$$c_{gg}^{\rm SM} \simeq \frac{1}{12\pi^2}\left[A_f\left(\frac{m_h^2}{4m_t^2}\right) + A_f\left(\frac{m_h^2}{4m_b^2}\right)\right],$$

$$A_f(\tau) \equiv \frac{3}{2\tau^2}\left[(\tau-1)f(\tau)+\tau\right],$$

$$f(\tau) \equiv \begin{cases} \arcsin^2\sqrt{\tau}, & \tau \le 1 \\ -\frac{1}{4}\left[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2, & \tau > 1 \end{cases}. \tag{4.3}$$

As discussed in ref. [88], in this case it is appropriate to calculate $c_{gg}^{\rm SM}$ at the leading order in QCD because then the large k-factors, approximately common for c_{gg} and δy_u, cancel in the ratio [88b]. Numerically,

$$\hat{c}_{gg} \simeq c_{gg} + (8.7\delta y_u - (0.3-0.3i)\delta y_d) \times 10^{-3},$$

$$c_{gg}^{\rm SM} \simeq (8.4+0.3i) \times 10^{-3}, \tag{4.4}$$

$$\frac{\sigma_{ggh}}{\sigma_{ggh}^{\rm SM}} \simeq 1 + 237c_{gg} + 2.06\delta y_u - 0.06\delta y_d. \tag{4.5}$$

(2) Vector boson fusion (VBF), $qq \to hqq$:

$$\frac{\sigma_{\rm VBF}}{\sigma_{\rm VBF}^{\rm SM}} \simeq 1 + 1.49\delta c_w + 0.51\delta c_z - \begin{pmatrix}1.08\\1.11\\1.23\end{pmatrix} c_{w\Box} - 0.10c_{ww} - \begin{pmatrix}0.35\\0.35\\0.40\end{pmatrix} c_{z\Box} - 0.04c_{zz} - 0.10c_{\gamma\Box} - 0.02c_{z\gamma}$$
$$\to 1 + 2\delta c_z - 2.25c_{z\Box} - 0.83c_{zz} + 0.30c_{z\gamma} + 0.12c_{\gamma\gamma}. \tag{4.6}$$

The numbers in the columns multiplying $c_{w\Box}$ and $c_{z\Box}$ refer to the LHC collision energy of $\sqrt{s} = 7$, 8, and 13 TeV; for other parameters the dependence is weaker. The expression after the arrow arises due to replacement of the dependent couplings by the independent ones in eq. (3.2). Each LHC Higgs analysis uses somewhat different cuts to isolate the VBF signal, and the relative cross-section slightly depends on these cuts. The result in eq. (4) has been computed numerically by simulating the parton-level process at the tree level in MadGraph5 [90] using the default `nn23lo1` PDF and the cuts $p_{T,q} > 20$ GeV, $|\eta_q| < 5$ and $m_{qq} > 250$ GeV. Replacing the last cut by $m_{qq} > 500$ GeV affects the numbers at the level of 5%.

(3) Vector boson associated production (Vh), $q\bar{q} \to Vh$, where $V = W, Z$,

$$\frac{\sigma_{Wh}}{\sigma_{Wh}^{\rm SM}} \simeq 1+2\delta c_w + \begin{pmatrix}6.39\\6.51\\6.96\end{pmatrix} c_{w\Box} + \begin{pmatrix}1.49\\1.49\\1.50\end{pmatrix} c_{ww}$$
$$\to 1+2\delta c_z + \begin{pmatrix}9.26\\9.43\\10.08\end{pmatrix} c_{z\Box} + \begin{pmatrix}4.35\\4.41\\4.63\end{pmatrix} c_{zz}$$
$$- \begin{pmatrix}0.81\\0.84\\0.93\end{pmatrix} c_{z\gamma} - \begin{pmatrix}0.43\\0.44\\0.48\end{pmatrix} c_{\gamma\gamma}$$
$$\frac{\sigma_{Zh}}{\sigma_{Zh}^{\rm SM}} \simeq 1+2\delta c_z + \begin{pmatrix}5.30\\5.40[2pt]\\5.72\end{pmatrix} c_{z\Box} + \begin{pmatrix}1.79\\1.80\\1.82\end{pmatrix} c_{zz}$$
$$+ \begin{pmatrix}0.80\\0.82\\0.87\end{pmatrix} c_{\gamma\Box} + \begin{pmatrix}0.22\\0.22\\0.22\end{pmatrix} c_{z\gamma},$$
$$\to 1+2\delta c_z + \begin{pmatrix}7.61\\7.77\\8.24\end{pmatrix} c_{z\Box} + \begin{pmatrix}3.31\\3.35\\3.47\end{pmatrix} c_{zz}$$
$$- \begin{pmatrix}0.58\\0.60\\0.65\end{pmatrix} c_{z\gamma} + \begin{pmatrix}0.27\\0.28\\0.30\end{pmatrix} c_{\gamma\gamma}. \tag{4.7}$$

The numbers in the columns refer to the LHC collision energy of $\sqrt{s} = 7$, 8, and 13 TeV.

(4) Top pair associated production, $gg \to ht\bar{t}$:

$$\frac{\sigma_{tth}}{\sigma_{tth}^{\rm SM}} \simeq 1 + 2\delta y_u. \tag{4.8}$$

Decay

(1) $h \to f\bar{f}$. Higgs boson decays into 2 fermions occur at the tree level in the SM via the Yukawa couplings in eq. (2.17). In the presence of $D = 6$ operators they are affected via the corrections to the Yukawa couplings in eq. (3.5):

$$\frac{\Gamma_{cc}}{\Gamma_{cc}^{\rm SM}} \simeq 1 + 2\delta y_u, \quad \frac{\Gamma_{bb}}{\Gamma_{bb}^{\rm SM}} \simeq 1 + 2\delta y_d,$$
$$\frac{\Gamma_{\tau\tau}}{\Gamma_{\tau\tau}^{\rm SM}} \simeq 1 + 2\delta y_e, \tag{4.9}$$

where $\Gamma(h \to Y) \equiv \Gamma_Y$.

(2) $h \to VV$. In the SM, Higgs decays into on-shell gauge bosons: gluon pairs gg, photon pairs $\gamma\gamma$, and $Z\gamma$ occur only at the one-loop level. In the presence of $D = 6$ operators these decays are corrected already at the tree level by the 2-derivative contact interactions of the Higgs boson with two vector bosons in eq. (3.3). The relative decay widths are given by

$$\frac{\Gamma_{VV}}{\Gamma_{VV}^{\rm SM}} \simeq \left|1 + \frac{\hat{c}_{vv}}{c_{vv}^{\rm SM}}\right|^2, \quad vv \in \{gg, \gamma\gamma, z\gamma\}, \tag{4.10}$$

where

$$\hat{c}_{\gamma\gamma} = c_{\gamma\gamma}, \quad c_{\gamma\gamma}^{\rm SM} \simeq -8.3 \times 10^{-2},$$
$$\hat{c}_{z\gamma} = c_{z\gamma}, \quad c_{z\gamma}^{\rm SM} \simeq -5.9 \times 10^{-2}, \tag{4.11}$$

while $\hat{c}_{gg}$ and $c_{gg}^{\rm SM}$ are defined in eq. (4.3). Note that contributions to $\Gamma_{\gamma\gamma}$ and $\Gamma_{z\gamma}$ arising due to corrections to the SM Higgs couplings to the W bosons and fermions are not included in eq. (4.11), unlike in eq. (4.3). The reason is that, for these processes, corrections from $D = 6$ operators are included at the tree level only. If these particular one-loop corrections were included, one should also consistently include *all* one-loop corrections to this process arising at the $D = 6$ level, some of which are divergent and require renormalization. The net result would be to redefine $\hat{c}_{\gamma\gamma} = c_{\gamma\gamma}^{\rm ren.} - 0.11\delta c_w + 0.02\delta y_u + \cdots$ and $\hat{c}_{z\gamma} = c_{z\gamma}^{\rm ren.} - 0.06\delta c_w + 0.003\delta y_t + \cdots$. Here 'ren.' stands for 'renormalized' and the dots stand for a dependence on other Lagrangian parameters (c_{ww}, $c_{w\Box}$, and corrections to triple gauge couplings). A full next-to-leading order computation of these processes have not been yet attempted in the literature.

(3) $h \to 4f$. The decay process $h \to 2\ell 2\nu$ (where ℓ here stands for charged leptons) proceeds via intermediate W bosons. The relative width is given by

$$\frac{\Gamma_{2\ell 2\nu}}{\Gamma_{2\ell 2\nu}^{\rm SM}} \simeq 1 + 2\delta c_w + 0.46 c_{w\Box} - 0.15 c_{ww}$$
$$\to 1 + 2\delta c_z + 0.67 c_{z\Box} + 0.05 c_{zz}$$
$$- 0.17 c_{z\gamma} - 0.05 c_{\gamma\gamma}. \tag{4.12}$$

In the SM, the decay process $h \to 4\ell$ proceeds at the tree level via intermediate Z bosons. In the presence of $D = 6$ operators, intermediate photon contributions may also arise at the tree level. If that is the case, the decay width diverges due to the photon pole. Below, the relative width $\bar{\Gamma}(h \to 4\ell)$ regulated by imposing the cut $m_{\ell\ell} > 12$ GeV on the invariant mass of same-flavour lepton pairs is quoted:

$$\frac{\bar{\Gamma}_{4\ell}}{\bar{\Gamma}^{\rm SM}_{4\ell}} \simeq 1+2\delta c_z+\begin{pmatrix}0.41\\0.39\end{pmatrix}c_{z\Box}-\begin{pmatrix}0.15\\0.14\end{pmatrix}c_{zz}$$
$$+\begin{pmatrix}0.07\\0.05\end{pmatrix}c_{z\gamma}-\begin{pmatrix}0.02\\0.02\end{pmatrix}c_{\gamma\Box}+\begin{pmatrix}<0.01\\0.03\end{pmatrix}c_{\gamma\gamma}$$
$$\to 1+2\delta c_z+\begin{pmatrix}0.35\\0.32\end{pmatrix}c_{z\Box}-\begin{pmatrix}0.19\\0.19\end{pmatrix}c_{zz}$$
$$+\begin{pmatrix}0.09\\0.08\end{pmatrix}c_{z\gamma}+\begin{pmatrix}0.01\\0.02\end{pmatrix}c_{\gamma\gamma}. \qquad (4.13)$$

The numbers in the columns correspond to the $2e2\mu$ and $4e/\mu$ final states, respectively. The difference between these two is numerically irrelevant in the total width, but may be important for differential distributions, especially regarding the $c_{\gamma\gamma}$ dependence [91]. The dependence on the $m_{\ell\ell}$ cut is weak; very similar numbers are obtained if $m_{\ell\ell} > 4$ GeV is imposed instead.

Given the partial widths, the branching fractions can be computed as $\mathrm{Br}_Y = \Gamma_Y/\Gamma(h \to \mathrm{all})$, where the total decay width is given by

$$\frac{\Gamma(h\to\mathrm{all})}{\Gamma(h\to\mathrm{all})} \simeq \frac{\Gamma_{bb}}{\Gamma^{\rm SM}_{bb}}\mathrm{Br}^{\rm SM}_{bb}+\frac{\Gamma_{cc}}{\Gamma^{\rm SM}_{cc}}\mathrm{Br}^{\rm SM}_{cc}$$
$$+\frac{\Gamma_{\tau\tau}}{\Gamma^{\rm SM}_{\tau\tau}}\mathrm{Br}^{\rm SM}_{\tau\tau}+\frac{\Gamma_{WW^*}}{\Gamma^{\rm SM}_{WW^*}}\mathrm{Br}^{\rm SM}_{WW^*}$$
$$+\frac{\Gamma_{ZZ^*}}{\Gamma^{\rm SM}_{ZZ^*}}\mathrm{Br}^{\rm SM}_{ZZ^*}+\frac{\Gamma_{gg}}{\Gamma^{\rm SM}_{gg}}\mathrm{Br}^{\rm SM}_{gg}. \qquad (4.14)$$

Note that, in line with the basic assumption of no new light particles, there is no additional contributions to the Higgs width other than from the SM decay channels. In particular, the invisible Higgs width is absent in this EFT framework (except for the small SM contribution arising via $h \to ZZ^* \to 4\nu$) [91a].

5. Current constraints

This section presents the constraints on the independent couplings characterizing the Higgs boson couplings in the dimension-six EFT Lagrangian. A disclaimer is in order. The objective is to illustrate what is the constraining power of the present data. As we shall see, the existing data is not yet good enough to even constrain all the couplings inside the EFT validity range. In the future, as the measurements become more precise and more information is available, this kind of analysis will become fully consistent.

5.1 *Data*

Here, the experimental data used to constrain the effective theory parameters are reviewed first. In the best of all worlds, the LHC Collaborations would quote a multidimensional likelihood function for the signal strength $\mu_{X;Y}$ for all production modes and decay channels, separately for each LHC collision energy. This would allow one to consistently use available experimental information, including non-trivial correlations between the different μs. Although the manner

Table 2. The LHC Higgs results used in the fit (see §5.1 for explanations).

ATLAS				CMS			
Channel	μ	Production	Ref.	Channel	μ	Production	Ref.
$\gamma\gamma$	$1.17^{+0.28}_{-0.26}$	cats.	[92]	$\gamma\gamma$	$1.12^{+0.25}_{-0.22}$	cats.	[101]
$Z\gamma$	$2.7^{+4.6}_{-4.5}$	total	[93]	$Z\gamma$	$-0.2^{+4.9}_{-4.9}$	total	[102]
ZZ^*	$1.46^{+0.40}_{-0.34}$	2D	[94]	ZZ^*	$1.00^{+0.29}_{-0.29}$	2D	[103]
WW^*	$1.18^{+0.24}_{-0.21}$	2D	[95]	WW^*	$0.83^{+0.21}_{-0.21}$	2D	[103]
	$2.1^{+1.9}_{-1.6}$	Wh	[96]		$0.80^{+1.09}_{-0.93}$	Vh	[103]
	$5.1^{+4.3}_{-3.1}$	Zh	[96]	$\tau\tau$	$0.91^{+0.28}_{-0.28}$	2D	[103]
$\tau\tau$	$1.44^{+0.42}_{-0.37}$	2D	[97]		$0.87^{+1.00}_{-0.88}$	Vh	[103]
bb	$1.11^{+0.65}_{-0.61}$	Wh	[98]		$-1.3^{+6.3}_{-5.5}$	tth	[104]
	$0.05^{+0.52}_{-0.49}$	Zh	[98]	bb	$0.89^{+0.47}_{-0.44}$	Vh	[103]
	$1.5^{+1.1}_{-1.1}$	tth	[99]		$1.2^{+1.6}_{-1.5}$	tth	[105]
$\mu\mu$	$-0.7^{+3.7}_{-3.7}$	total	[93]	$\mu\mu$	$0.8^{+3.5}_{-3.4}$	total	[106]
multi-ℓ	$2.1^{+1.4}_{-1.2}$	tth	[100]	multi-ℓ	$3.8^{+1.4}_{-1.4}$	tth	[104]

in which the LHC data are presented has been constantly improving, we are not yet in the ideal world. For these reasons, constraining Higgs couplings from the existing data involves inevitably somewhat arbitrary assumptions and approximations. Nevertheless, thanks to the fact that the experimental uncertainties are still statistics-dominated in most cases, one should expect that these approximations do not affect the results in a dramatic way.

The measurements of the Higgs signal strength μ included in this analysis are summarized in table 2. They are separated according to the final state (channel) and the production mode. For the all-inclusive production mode (total), I use the value of μ quoted in table 2 [91b]. The same is true for μ in a specific production mode (*Wh, Zh, tth*), in which case correlations with other production channels are ignored. In the remaining cases, μ is quoted for illustration only, and more information is included in the analysis. 2D stands for two-dimensional likelihood functions in the plane $\mu_{ggh+tth}$–$\mu_{\mathrm{VBF}+Vh}$. As the contribution of the Vh production mode is subleading with respect to the VBF one, separate measurements of the Vh signal strength are combined (whenever it is given) with the 2D likelihood, ignoring the correlation between the two. For the diphoton final state, five-dimensional likelihood function is constructed in the space spanned by $(\mu_{ggh}, \mu_{tth}, \mu_{\mathrm{VBF}}, \mu_{Wh}, \mu_{Zh})$ using the signal strength in all diphoton event categories (cats.), using the known contribution of each production mode to each category. In many channels, there is a certain degree of arbitrariness as to which set of results (inclusive, 1D, 2D, or cats.) to include in the fit; here the strategy is to choose the set that maximizes the available information about various EFT couplings.

5.2 *Fit*

Using the dependence of the signal strength on EFT parameters worked out in §4 and the LHC data in table 2, one can constrain all CP-even independent Higgs couplings in eq. (3.2) [106a]. In the Gaussian approximation near the best-fit point, the following constraints can be seen:

$$\begin{pmatrix} \delta c_z \\ c_{zz} \\ c_{z\Box} \\ c_{\gamma\gamma} \\ c_{z\gamma} \\ c_{gg} \\ \delta y_u \\ \delta y_d \\ \delta y_e \end{pmatrix} = \begin{pmatrix} -0.12 \pm 0.20 \\ 0.5 \pm 1.8 \\ -0.21 \pm 0.82 \\ 0.014 \pm 0.029 \\ 0.01 \pm 0.10 \\ -0.0056 \pm 0.0028 \\ 0.55 \pm 0.30 \\ -0.42 \pm 0.45 \\ -0.17 \pm 0.35 \end{pmatrix}, \tag{5.1}$$

where the uncertainties correspond to 1σ. The correlation matrix is

$$\rho = \begin{pmatrix} 1 & -0.23 & 0.17 & -0.62 & -0.18 & 0.16 & 0.09 & 0.88 & 0.63 \\ \cdot & 1 & -0.997 & 0.85 & 0.23 & 0.13 & 0.17 & -0.47 & -0.81 \\ \cdot & \cdot & 1 & -0.82 & -0.21 & -0.15 & -0.17 & 0.41 & 0.78 \\ \cdot & \cdot & \cdot & 1 & 0.27 & 0.02 & 0.09 & -0.79 & -0.92 \\ \cdot & \cdot & \cdot & \cdot & 1 & 0.01 & 0.02 & -0.22 & -0.26 \\ \cdot & \cdot & \cdot & \cdot & \cdot & 1 & -0.81 & 0.21 & 0.03 \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & 1 & 0.05 & -0.06 \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & 1 & 0.82 \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & 1 \end{pmatrix}. \tag{5.2}$$

Using the above central values c_0, uncertainties δc, and the correlation matrix ρ, one can reconstruct the nine-dimensional likelihood function near the best-fit point:

$$\chi^2 \simeq \sum_{ij} [c - c_0]_i \sigma_{ij}^{-2} [c - c_0]_j,$$
$$\sigma_{ij}^{-2} \equiv [[\delta c]_i \rho_{ij} [\delta c]_j]^{-1}. \tag{5.3}$$

5.3 *Discussion*

As one can see from eq. (5.1), certain EFT parameters are very weakly constrained by experiment, with order one deviations from the SM being allowed. In other words, the current data cannot even constrain all the parameters to be within the EFT validity range. This violates the initial assumption that the $D = 6$ operators give a small correction on top of the SM. For this reason, the results in eq. (5.1) should not be taken at face value. In particular, one should conclude that there is currently no model-independent constraints at all on c_{zz} and $c_{z\Box}$. Indeed, including corrections to observables that are quadratic in these parameters would completely change the central values and the uncertainties. This signals a sensitivity of the fit to operators with $D > 6$. Furthermore, the experimental constraints in the $Z\gamma$ channel are still too weak to justify the linear approximation. Again, including quadratic EFT corrections would significantly affect the constraints on $c_{z\gamma}$. To a lesser extent, the sensitivity to higher-order EFT corrections is also true for deformations along the δy_f directions.

Nevertheless, the results in eq. (5.1) are of some value. First of all, they demonstrate that certain EFT parameters are strongly constrained. This is true especially for c_{gg} and $c_{\gamma\gamma}$ which are constrained at the 10^{-3} level. Next, the fit in eq. (5.1) identifies 'blind' directions in the space of the EFT parameters that are weakly constrained by current data. The most dramatic example is the approximate degeneracy along the line $c_{zz} \approx -2.3 c_{z\Box}$, as witnessed by the ≈ 1 entry in the correlation matrix in eq. (5.2). More data are needed to lift this degeneracy. To this end, extremely helpful pieces of information can be extracted from differential distributions in $h \to 4\ell$ decays [108,115–119], as well as in the *Vh* [42,120–125] and VBF [126–128] production. A consistent, model-independent EFT approach to Higgs differential distributions has not yet been implemented in LHC analyses, but the CMS Collaboration made first steps in this direction [129]. Note also the large correlations between δy_d and other parameters. This happens because δy_d strongly affects the total Higgs width (via the $h \to b\bar{b}$ partial width) and this way it affects the signal strength in all Higgs decay channels. More precise measurements of the signal strength in the $h \to b\bar{b}$ channel should soon alleviate this degeneracy. Finally, there is the strong correlation between c_{gg} and δy_t which has been extensively discussed in the Higgs fits literature. In the future, that degeneracy will be lifted by better measurements of the *tth* signal strength, and by the measurements of the Higgs p_{T} distribution in the gluon fusion production mode [130–134] (see also [135] for an earlier work in this direction).

Finally, the importance of the fit is in the fact that the likelihood in eq. (5.3) can be combined with other datasets that constrain the same EFT parameters. In this case, one may obtain stronger bounds that will push the parameters into the EFT validity range. For example, one can use constraints on cubic self-couplings of electroweak gauge bosons [83,86,119,136–139]. These are customarily parametrized by three parameters $\delta g_{1,z}$, $\delta\kappa_\gamma$, λ_z [140] which characterize deviations of these self-couplings from the SM predictions. Now, in the EFT Lagrangian with $D = 6$ operators, the first two parameters are related to the Higgs couplings. In the Higgs basis one finds [26]

$$\delta g_{1,z} = \frac{1}{2(g_L^2 - g_Y^2)}\left[c_{\gamma\gamma} e^2 g_Y^2 + c_{z\gamma}(g_L^2 - g_Y^2) g_Y^2 - c_{zz}(g_L^2 + g_Y^2) g_Y^2 - c_{z\Box}(g_L^2 + g_Y^2) g_L^2\right],$$
$$\delta\kappa_\gamma = -\frac{g_L^2}{2}\left(c_{\gamma\gamma}\frac{e^2}{g_L^2 + g_Y^2} + c_{z\gamma}\frac{g_L^2 - g_Y^2}{g_L^2 + g_Y^2} - c_{zz}\right). \tag{5.4}$$

Therefore, model-independent constraints on triple gauge couplings imply additional constraint on the EFT parameters characterizing the Higgs couplings. In particular, Falkowski and Riva [138] argue that, after marginalizing over λ_z, the single and pair W boson production in LEP-2 implies the bounds $\delta g_{1,z} = -0.83 \pm 0.34$, $\delta\kappa_\gamma = 0.14 \pm 0.05$ with the correlation coefficient $[\rho]_{\delta g_{1,z}}^{\delta\kappa_\gamma} = -0.71$. Combining this bound with the likelihood in eq. (5.3) the degeneracy between c_{zz} and $c_{z\Box}$ is lifted, and one obtains much stronger bounds: $c_{zz} = 0.22 \pm 0.18$, $c_{z\Box} = -0.08 \pm 0.09$, $[\rho]_{c_{zz}}^{c_{z\Box}} = -0.76$. More constraints of this type, for example model-independent constraints on triple gauge couplings from the LHC, could further improve the limits on Higgs couplings within the EFT approach. As soon as more precise Higgs and diboson data from the 13 TeV LHC run start arriving, it should be possible to constrain all the nine parameters in eq. (5.1) safely within the EFT validity range.

6. Closing words

The Higgs boson has been discovered, and for the remainder of this century we shall study its properties. Precision measurements of Higgs couplings and determination of their tensor structure is an important part of the physics programme at the LHC and future colliders. Given that not the slightest hint for a particular scenario beyond the SM has emerged so far, it is important to (also) perform these studies in a model-independent framework. The EFT approach described here, with the SM extended by dimension-six operators, provides a perfect tool to this end.

One should be aware that Higgs precision measurements cannot probe new physics at very high scales. For example, LHC Higgs measurements are sensitive to new physics at $\Lambda \sim 1$ TeV at the most. This is not too impressive, especially compared to the new physics reach of flavour observables or even electroweak precision tests. However, Higgs physics probes a subset of operators that are often not accessible by other searches. For example, for most of the nine parameters in eq. (5.1) the only experimental constraints come from Higgs physics. It is certainly conceivable that new physics talks to the SM via the Higgs portal, and it will first manifest itself within this particular class of $D = 6$ operators. If this is the case, we must not miss it.

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[9a] One could consider a more general EFT where assumptions 2 and 3 are also relaxed and the electroweak symmetry is realized non-linearly [10–13]. In that case, the Higgs boson is introduced as a singlet of the local symmetry, rather than as a part of an $SU(2)$ doublet.
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while contributions of higher-order operators are suppressed by more powers of Λ. One example is the lepton-flavour violating Higgs decays into two fermions where the SM contribution is exactly zero. In this review, I focus on the observables where the SM contribution is dominant.

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[88b] Accidentally, with the SM parameters used in this review, the dependence on δy_d is also captured with a decent accuracy by this procedure. One can compare eq. (4.5) to NLO QCD results in ref. [89], where the coefficient in front of δy_d is found to be −0.06 for $\sqrt{s} = 8$ TeV and −0.05 for $\sqrt{s} =$14 TeV.

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[91b] CMS does not quote the best-fit μ in the $Z\gamma$ channel. The value in table 2 was obtained by digitizing the plot showing the expected and observed 95% CL limits on μ as a function of m_h, extracting the values at $m_h = 125$ GeV, and using these to calculate the best-fit μ assuming that the uncertainties are Gaussian. This is a dire reminder of how Higgs fits had to be done back in the early 2010s.

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