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Handbook of LHC Higgs cross sections:
4. Deciphering the nature of the Higgs sector

Report of the LHC Higgs Cross Section Working Group

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Abstract

This Report summarises the results of the activities in the period 2014–2015 of the LHC Higgs Cross Section Working Group. The main goal of the working group was to present the state of the art of Higgs Physics at the LHC, integrating all new results that have appeared in the last few years. This report follows the first working group report *Handbook of LHC Higgs Cross Sections: 1. Inclusive Observables* (CERN-2011-002) and the second working group report *Handbook of LHC Higgs Cross Sections: 2. Differential Distributions* (CERN-2012-002). *Handbook of LHC Higgs Cross Sections: 3. Higgs properties* (CERN-2013-004). *BlaBla...*

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Chapter 1

EFT formalism

For a large class of models beyond the SM, physics at energies below the mass scale Λ of the new particles can be parametrized by an effective field theory (EFT) where the SM Lagrangian is supplemented by new operators with canonical dimensions D larger than 4. The theory has the same field content and the same linearly realized $SU(3) \times SU(2) \times U(1)$ local symmetry as the SM.[1] The higher-dimensional operators are organized in a systematic expansion in D, where each consecutive term is suppressed by a larger power of Λ. For a general introduction to the EFT formalism see e.g. [1–4]; for recent review articles about EFT in connection with Higgs physics see e.g. [5–9].

Quite generally, the EFT Lagrangian takes the form:

$$\mathcal{L}_{\rm eff} = \mathcal{L}_{\rm SM} + \sum_i \frac{c_i^{(5)}}{\Lambda}\mathcal{O}_i^{(5)} + \sum_i \frac{c_i^{(6)}}{\Lambda^2}\mathcal{O}_i^{(6)} + \sum_i \frac{c_i^{(7)}}{\Lambda^3}\mathcal{O}_i^{(7)} + \sum_i \frac{c_i^{(8)}}{\Lambda^4}\mathcal{O}_i^{(8)} + \cdots, \tag{1.1}$$

where each $\mathcal{O}_i^{(D)}$ is an $SU(3) \times SU(2) \times U(1)$ invariant operator of dimension D and the parameters $c_i^{(D)}$ multiplying the operators in the Lagrangian are called the *Wilson coefficients*. This EFT is intended to parametrize observable effects of a large class of BSM theories where new particles, with mass of order Λ, are much heavier than the SM ones and much heavier than the energy scale at which the experiment is performed. The main motivation to use this framework is that the constraints on the EFT parameters can be later re-interpreted as constraints on masses and couplings of new particles in many BSM theories. In other words, translation of experimental data into a theoretical framework has to be done only once in the EFT context, rather than for each BSM model separately.

The contribution of each $O_i^{(D)}$ to amplitudes of physical processes at the energy scale of order v scales[2] as $(v/\Lambda)^{D-4}$. Since $v/\Lambda < 1$ by construction, the EFT in its validity regime typically describes *small* deviations from the SM predictions, although, under certain conditions discussed later in this chapter, it may be consistent to use this framework to describe large deviations.

A complete and non-redundant set of operators that can be constructed from the SM fields is known for D=5 [10], D=6 [11], D=7 [12], and D=8 [13]. All D=5 operators violate the lepton number [10], while all D=7 operators violate $B - L$ [12] (the latter is true for all odd-D operators [14]). Then, experimental constraints dictate that their Wilson coefficients must be suppressed at a level which makes them unobservable at the LHC [15], and for this reason D=5 and 7 operators will not be discussed here. Consequently, the leading new physics effects are expected from operators with D=6 [16], whose contributions scale as $(v/\Lambda)^2$. Contributions from operators with $D \geq 8$ are suppressed by at least $(v/\Lambda)^4$, and in most of the following discussion we will assume that they can be neglected.

In the rest of this chapter, we discuss in detail the set D=6 operators that can be constructed from the SM fields. We review various possible choices of these operators (the so-called *basis*) and their phenomenological effects. We also discuss the validity regime of the SM EFT with D=6 operators. Only the operators that conserve the baryon and lepton numbers are considered. On the other hand, we do not impose a-priori any flavor symmetry. Also, we include CP violating operators in our discussion.

In Section 1 we introduce the SM Lagrangian extended by dimension-6 operators. Two popular bases of dimension-6 operators using the manifestly $SU(2) \times U(1)$ invariant formalism are explicitly listed. In Section 2 we discuss the interactions of the SM mass eigenstates that arise in the presence of dimension-6 operators beyond the SM, with the emphasis on the Higgs interactions. We also derive provide a map between the couplings in that effective Lagrangian and

[1] The latter assumption can be relaxed, leading to an EFT with a non-linearly realized electroweak symmetry. This framework is discussed in Section **??**.

[2] Apart from the scaling with Λ, the effects of higher-dimensional operators also scale with appropriate powers of couplings in the UV theory. The latter is important to assess the validity range of the EFT description, as discussed in Section 4.

Wilson coefficients of dimension-6 operators introduced in Section 1. In Section 3 we define a new basis of D=6 operators, the so-called Higgs basis, which is spanned by a subset of the independent couplings of the mass eigenstate Lagrangian. In Section 4 we discuss under which conditions and in which energy range does the EFT with D=6 operators provides an adequate description of the underlying theory beyond the SM. We also comment on several physically important examples where such a framework is insufficient. This chapter attempts to review the most important results from the point of view of LHC Higgs phenomenology. Some additional details and derivations can be found in the associated LHCHXSWG internal note [17].

1 SM EFT with dimension-6 operators

We consider an EFT Lagrangian where the SM is extended by dimension-6 operators:

$$\mathcal{L}_{\rm EFT} = \mathcal{L}_{\rm SM} + \frac{1}{\Lambda^2}\sum_i c_i^{(6)} O_i^{(6)}. \tag{1.2}$$

To fix our notation and conventions, we first write down the SM Lagrangian:

$$\begin{aligned}
\mathcal{L}_{\rm SM} &= -\frac{1}{4}G^a_{\mu\nu}G^a_{\mu\nu} - \frac{1}{4}W^i_{\mu\nu}W^i_{\mu\nu} - \frac{1}{4}B_{\mu\nu}B_{\mu\nu} + D_\mu H^\dagger D_\mu H + \mu_H^2 H^\dagger H - \lambda(H^\dagger H)^2 \\
&+ \sum_{f\in q,\ell} i\bar{f}_L\gamma_\mu D_\mu f_L + \sum_{f\in u,d,e} i\bar{f}_R\gamma_\mu D_\mu f_R \\
&- \left[\tilde{H}^\dagger \bar{u}_R y_u q_L + H^\dagger \bar{d}_R y_d V_{\rm CKM}^\dagger q_L + H^\dagger \bar{e}_R y_e \ell_L + {\rm h.c.}\right].
\end{aligned} \tag{1.3}$$

Here, G^a_μ, W^i_μ, and B_μ denote the gauge fields of the $SU(3)\times SU(2)\times U(1)$ local symmetry. The corresponding gauge couplings are denoted by g_s, g, g'; we also define the electromagnetic coupling $e = gg'/\sqrt{g^2+g'^2}$, and the Weinberg angle $s_\theta = g'/\sqrt{g^2+g'^2}$. The field strength tensors are defined as $G^a_{\mu\nu} = \partial_\mu G^a_\nu - \partial_\nu G^a_\mu + g_s f^{abc} G^b_\mu G^c_\nu$, $W^i_{\mu\nu} = \partial_\mu W^i_\nu - \partial_\nu W^i_\mu + g\epsilon^{ijk}W^j_\mu W^k_\nu$, $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$. The Higgs doublet is denoted as H, and we also define $\tilde{H}_i = \epsilon_{ij}H^*_j$. It acquires the VEV $\langle H^\dagger H\rangle = v^2/2$. In the unitary gauge we have $H = (0,(v+h)/\sqrt{2})$, where h is the Higgs boson field. After electroweak symmetry breaking, the electroweak gauge boson mass eigenstates are defined as $W^\pm = (W^1 \mp iW^2)/\sqrt{2}$, $Z = c_\theta W^3 - s_\theta B$, $A = s_\theta W^3 + c_\theta B$, where $c_\theta = \sqrt{1-s_\theta^2}$. The tree-level masses of W and Z bosons are given by $m_W = gv/2$, $m_Z = \sqrt{g^2+g'^2}v/2$. The left-handed Dirac fermions $q_L = (u_L, V_{\rm CKM}d_L)$ and $\ell_L = (\nu_L, e_L)$ are doublets of the SU(2) gauge group, and the right-handed Dirac fermions u_R, d_R, e_R are SU(2) singlets. All fermions are 3-component vectors in the generation space, and y_f are 3×3 matrices. We work in the basis where the fermion mass matrix is diagonal with real, positive entries. In this basis, y_f are diagonal, and the fermion masses are given by $m_{f_i} = v[y_f]_{ii}/\sqrt{2}$. The 3 electroweak parameters g, g', v are customarily derived from the Fermi constant G_F measured in muon decays, Z boson mass m_Z, and the low-energy electromagnetic coupling $\alpha(0)$. The tree-level relations between the input observables and the electroweak parameters are given by:

$$G_F = \frac{1}{\sqrt{2}v^2}, \qquad \alpha = \frac{g^2 g'^2}{4\pi(g^2+g'^2)}, \qquad m_Z = \frac{\sqrt{g^2+g'^2}v}{2}. \tag{1.4}$$

We demand that the dimension-6 operators $O_i^{(6)}$ in Eq. (1.2) form a complete, non-redundant set - a so-called *basis*. Complete means that any dimension-6 operator is either a part of the basis or can be obtained from a combination of operators in the basis using equations of motion, integration by parts, field redefinitions, and Fierz transformations. Non-redundant means it is a minimal such set. Any complete basis leads to the same physical predictions concerning possible new physics effects. Several bases have been proposed in the literature, and they may be convenient for specific applications. Below we describe two popular choices in the existing literature. Later, in Section 3, we propose a new basis choice that is particularly convenient for leading-order LHC Higgs analyses in the EFT framework.

Historically, a complete and non-redundant set of D=6 operators was first identified in Ref. [11], and is usually referred to as the *Warsaw basis*. For our purpose, it is more convenient to work with a variant of that basis which differs from the one in Ref. [11] by the following aspects:

- We replace the operator $|H^\dagger D_\mu H|^2$ by $O_T = (H^\dagger \overleftrightarrow{D_\mu} H)^2$, where $H^\dagger \overleftrightarrow{D_\mu} H \equiv H^\dagger D_\mu H - D_\mu H^\dagger H$. The reason is that O_T is more directly connected to violation of custodial symmetry among Higgs couplings.
- For Yukawa-type D=6 operators $H|H|^2\bar{f}f$ we subtracted v^2 from $|H|^2$in the definition, so that they do not contribute to fermion mass terms. This way we avoid tedious rotations of the fermion fields to bring them back to the mass eigenstate basis. Moreover, we isolated factor of fermion masses in the definition, for a more direct connection to minimal flavor violating scenarios. Starting with the Yukawa couplings $-H\bar{f}'_R(Y'_f + c'_f H^\dagger H/v^2)f'_L$ we can

bring them to the form in Eq. (1.3) and Table 1.2 by defining $f'_{L,R} = U_{L,R} f_{L,R}$, $\sqrt{m_i m_j}[c_f]_{ij}/v = [U_R^\dagger c'_f U_L]_{ij}$, $Y_f = U_R^\dagger (Y'_f + c'_f/2) U_L$, where $U_{L,R}$ are unitary rotations to the mass eigenstate basis.

For other operators, we often use a different notation and normalizations than the original reference. The full set of operators in the Warsaw basis is given in Tables 1.1, 1.2, and 1.3.

Another D=6 basis choice commonly used in the literature is the SILH basis [18, 19]. In the case we also use a different notation and normalization than in the original references. Compared to the Warsaw basis, the SILH basis introduces the following nine new operators:

$$\begin{aligned}
O_W &= \frac{ig}{2}\left(H^\dagger \sigma^i \overleftrightarrow{D_\mu} H\right) D_\nu W^i_{\mu\nu}, \\
O_B &= \frac{ig'}{2}\left(H^\dagger \overleftrightarrow{D_\mu} H\right) \partial_\nu B_{\mu\nu}, \\
O_{HW} &= ig\left(D_\mu H^\dagger \sigma^i D_\nu H\right) W^i_{\mu\nu}, \\
O_{HB} &= ig'\left(D_\mu H^\dagger D_\nu H\right) B_{\mu\nu}, \\
O_{\widetilde{HW}} &= ig\left(D_\mu H^\dagger \sigma^i D_\nu H\right) \widetilde{W}^i_{\mu\nu}, \\
O_{\widetilde{HB}} &= ig'\left(D_\mu H^\dagger D_\nu H\right) \widetilde{B}_{\mu\nu}, \\
O_{2W} &= D_\mu W^i_{\mu\nu} D_\rho W^i_{\rho\nu}, \\
O_{2B} &= \partial_\mu B_{\mu\nu} \partial_\rho B_{\rho\nu}, \\
O_{2G} &= D_\mu G^a_{\mu\nu} D_\rho G^a_{\rho\nu}.
\end{aligned} \tag{1.5}$$

Consequently, in order to have a non-redundant set of operators, 9 operators present in the Warsaw basis must be absent in the SILH basis. The absent ones are 4 bosonic operators O_{WW}, $O_{\widetilde{WW}}$, O_{WB}, $O_{\widetilde{WB}}$, 2 vertex operators $[O_{H\ell}]_{11}$, $[O'_{H\ell}]_{11}$, and 3 four-fermion operators $[O_{\ell\ell}]_{1221}$, $[O_{\ell\ell}]_{1122}$, $[O'_{uu}]_{3333}$.[3] The remaining operators are the same as in the Warsaw basis, and we use the normalizations in Tables 1.1, 1.2, and 1.3. There exists a 1-to- linear map between the Wilson coefficients in the Warsaw and SILH bases. The dictionary is given in Ref. [17].

2 Effective Lagrangian of mass eigenstates

In Section 1 we introduced an EFT with the SM supplemented by D=6 operators, using a manifestly $SU(2) \times U(1)$ invariant notation. At that point, the connection between the new operators and phenomenology is not obvious. To relate to high-energy collider observables, it is more transparent to work with the degrees of freedom that are mass eigenstates after electroweak symmetry breaking (Higgs boson, W, Z, photon, etc.). In this section we relate the Wilson coefficients of dimension-6 operators to the parameters of the tree-level effective Lagrangian describing the interactions of the mass eigenstates.

We demand that the effective Lagrangian at tree level written in term of mass eigenstates has the following features:

- #1 All kinetic and mass terms are diagonal and canonically normalized.
- #2 Tree-level relations between the electroweak parameters and input observables are the same as the SM ones in Eq. (1.4).
- #3 The non-derivative photon and gluon interactions with fermions are the same as in the SM.
- #4 Two-derivative self-interactions of the Higgs boson are absent.
- #5 For each fermion pair, the coefficient of the vertex-like Higgs interaction terms $\left(2\frac{h}{v} + \frac{h^2}{v^2}\right) V_\mu \bar{f} \gamma_\mu f$ is equal to the vertex correction to the respective $V_\mu \bar{f} \gamma_\mu f$ interaction.

These conditions greatly simplify the relation between the parameters of the Lagrangian and collider observables. In general, dimension-6 operators can induce interaction terms that do not respect these features. However, the conditions #1-#5 can always be achieved, *without any loss of generality*, by using equations of motion, integrating by parts, and redefining the fields and couplings. The required set of transformations starting from the Warsaw basis is presented in Ref. [17]; an analogous procedure could be executed starting from the SILH basis.

We move to discussing the interactions in the effective Lagrangian conditions once #1-#5 are satisfied. We will focus on interaction terms that are most relevant for LHC phenomenology. To organize the presentation, we split the Lagrangian into the following parts,

$$\mathcal{L}_{\rm EFT} = \mathcal{L}_{\rm kinetic} + \mathcal{L}_{\rm aff} + \mathcal{L}_{\rm vertex} + \mathcal{L}_{\rm dipole} + \mathcal{L}_{\rm tgc} + \mathcal{L}_{\rm hff} + \mathcal{L}_{\rm hvv} + \mathcal{L}_{hvff} + \mathcal{L}_{hdvff} + L_{h,\rm self} + \mathcal{L}_{h^2} + \mathcal{L}_{\rm other}. \tag{1.1}$$

[3] Refs. [18, 19] do not specify flavor indices of the absent operators when general flavor structure of D=6 operators is introduced. Here, for concreteness, we made a particular though somewhat arbitrary choice of these indices.

Below we define each term in order of appearance. We also express the corrections to the SM interactions in $\mathcal{L}_{\rm EFT}$ via Wilson coefficients of D=6 operators in the Warsaw and SILH basis. These corrections start at $\mathcal{O}(1/\Lambda^2)$ in the EFT expansion, and will ignore all $\mathcal{O}(1/\Lambda^4)$ and higher contributions. Below we only present the Lagrangian in the unitary gauge when the Goldstone bosons eaten by W and Z are set to zero; see Ref. [17] for a generalization to the R_ξ gauge.

Kinetic Terms

By construction, the kinetic terms of the mass eigenstates are diagonal and canonically normalized:

$$\begin{aligned}\mathcal{L}_{\rm kin} &= -\frac{1}{2}W^+_{\mu\nu}W^-_{\mu\nu} - \frac{1}{4}Z_{\mu\nu}Z_{\mu\nu} - \frac{1}{4}A_{\mu\nu}A_{\mu\nu} - \frac{1}{4}G^a_{\mu\nu}G^a_{\mu\nu} \\ &+ \frac{g^2v^2}{4}(1+\delta m)^2 W^+_\mu W^-_\mu + \frac{(g^2+g'^2)v^2}{8}Z_\mu Z_\mu \\ &+ \frac{1}{2}\partial_\mu h\partial_\mu h - \lambda v^2 h^2 + i\sum_{f\in q,\ell,u,d,e}\bar{f}\left(\gamma_\mu\partial_\mu - m_f\right)f. \end{aligned} \tag{1.2}$$

Above, the parameter λ is defined by the tree-level relation $m_h^2 = 2\lambda v^2$. There is no correction to the Z boson mass terms, in accordance with the condition #2. With this convention, the corrections to the W boson mass cannot be in general redefined away, and are parametrized by δm. The relation between δm and the Wilson coefficients in the Warsaw and SILH bases is given by

$$\begin{aligned}\delta m &= \frac{1}{g^2-g'^2}\left[-g^2g'^2c_{WB} + g^2c_T - g'^2\delta v\right] \\ &= -\frac{g^2g'^2}{4(g^2-g'^2)}\left(s_W + s_B + s_{2W} + s_{2B} - \frac{4}{g'^2}s_T + \frac{2}{g^2}[s'_{H\ell}]_{22}\right). \end{aligned} \tag{1.3}$$

For the sake of clarity, here and in the following denote the Wilson coefficients in the Warsaw basis by c_i, and in the SILH basis by s_i. We also define $\delta v = ([c'_{H\ell}]_{11} + [c'_{H\ell}]_{22})/2 - [c_{\ell\ell}]_{1221}/4$.

Gauge boson interactions with fermions

By construction (condition #3), the non-derivative photon and gluon interactions with fermions are the same as in the SM:

$$\mathcal{L}_{aff} = eA_\mu\sum_{f\in u,d,e}\bar{f}\gamma_\mu Q_f f + g_s G^a_\mu\sum_{f\in u,d}\bar{f}\gamma_\mu T^a f. \tag{1.4}$$

The analogous interactions of the W and Z boson may in general be affected by dimension-6 operators:

$$\begin{aligned}\mathcal{L}_{\rm vertex} &= \frac{g}{\sqrt{2}}\left(W^+_\mu\bar{\nu}_L\gamma_\mu\left(I_3+\delta g_L^{W\ell}\right)e_L + W^+_\mu\bar{u}\gamma_\mu\left(I_3+\delta g_L^{Wq}\right)d_L + W^+_\mu\bar{u}_R\gamma_\mu\delta g_R^{Wq}d_R + \text{h.c.}\right) \\ &+ \sqrt{g^2+g'^2}Z_\mu\left[\sum_{f\in u,d,e,\nu}\bar{f}_L\left(T^3_f - s_\theta^2Q_f + \gamma_\mu\delta g_L^{Zf}\right)f_L + \sum_{f\in u,d,e}\bar{f}_R\left(-s_\theta^2Q_f + \gamma_\mu\delta g_R^{Zf}\right)f_R\right] \end{aligned} \tag{1.5}$$

Here, I_3 is the 3×3 identity matrix, and the *vertex corrections* δg are 3×3 Hermitian matrices in the generation space, except for δg_R^{Wq} which is a general 3×3 complex matrix. The vertex corrections to W and Z boson couplings to fermions are expressed by the Wilson coefficients in the Warsaw basis as

$$\begin{aligned}\delta g_L^{W\ell} &= c'_{H\ell} + f(1/2,0) - f(-1/2,-1), \\ \delta g_L^{Z\nu} &= \frac{1}{2}c'_{H\ell} - \frac{1}{2}c_{H\ell} + f(1/2,0), \\ \delta g_L^{Ze} &= -\frac{1}{2}c'_{H\ell} - \frac{1}{2}c_{H\ell} + f(-1/2,-1), \\ \delta g_R^{Ze} &= -\frac{1}{2}c_{He} + f(0,-1), \end{aligned} \tag{1.6}$$

$$\begin{aligned}\delta g_L^{Wq} &= \left(c'_{Hq} + f(1/2,2/3) - f(-1/2,-1/3)\right)V_{\rm CKM}, \\ \delta g_R^{Wq} &= -\frac{1}{2}c_{Hud}, \\ \delta g_L^{Zu} &= \frac{1}{2}c'_{Hq} - \frac{1}{2}c_{Hq} + f(1/2,2/3), \end{aligned}$$

$$
\begin{aligned}
\delta g_L^{Zd} &= -\frac{1}{2}V_{\rm CKM}^\dagger c'_{Hq}V_{\rm CKM} - \frac{1}{2}V_{\rm CKM}^\dagger c_{Hq}V_{\rm CKM} + f(-1/2,-1/3),\\
\delta g_R^{Zu} &= -\frac{1}{2}c_{Hu} + f(0,2/3),\\
\delta g_R^{Zd} &= -\frac{1}{2}c_{Hd} + f(0,-1/3),
\end{aligned} \tag{1.7}
$$

where

$$
f(T_f^3,Q_f) = \left[-Q_f c_{WB}\frac{g^2g'^2}{g^2-g'^2} + (c_T-\delta v)\left(T_f^3 + Q_f\frac{g'^2}{g^2-g'^2}\right)\right]. \tag{1.8}
$$

The analogous expression in the SILH basis read

$$
\begin{aligned}
\delta g_L^{Z\nu} &= \frac{1}{2}s'_{H\ell} - \frac{1}{2}s_{H\ell} + \hat f(1/2,0),\\
\delta g_L^{Ze} &= -\frac{1}{2}s'_{H\ell} - \frac{1}{2}s_{H\ell} + \hat f(-1/2,-1),\\
\delta g_R^{Ze} &= -\frac{1}{2}s_{He} + \hat f(0,-1),\\
\delta g_L^{Zu} &= \frac{1}{2}s'_{Hq} - \frac{1}{2}s_{Hq} + \hat f(1/2,2/3),\\
\delta g_L^{Zd} &= -\frac{1}{2}V_{\rm CKM}^\dagger s'_{Hq}V_{\rm CKM} - \frac{1}{2}V_{\rm CKM}^\dagger s_{Hq}V_{\rm CKM} + \hat f(-1/2,-1/3),\\
\delta g_R^{Zu} &= -\frac{1}{2}s_{Hu} + \hat f(0,2/3),\\
\delta g_R^{Zd} &= -\frac{1}{2}s_{Hd} + \hat f(0,-1/3),\\
\delta g_L^{W\ell} &= s'_{H\ell} + \hat f(1/2,0) - \hat f(-1/2,-1),\\
\delta g_L^{Wq} &= \left(s'_{Hq} + \hat f(1/2,2/3) - \hat f(-1/2,-1/3)\right)V_{\rm CKM},\\
\delta g_R^{Wq} &= -\frac{1}{2}s_{Hud},
\end{aligned} \tag{1.9}
$$

where

$$
\begin{aligned}
\hat f(T_f^3,Q_f) &\equiv \frac{1}{4}\left[g^2 s_{2W} + g'^2 s_{2B} + 4s_T - 2[s'_{H\ell}]_{22}\right]T_f^3\\
&+ \frac{g'^2}{4(g^2-g'^2)}\left[-(2g^2-g'^2)s_{2B} - g^2(s_{2W}+s_W+s_B) + 4s_T - 2[s'_{H\ell}]_{22}\right]Q_f.
\end{aligned} \tag{1.10}
$$

Another type of gauge boson interactions with fermions are the so-called dipole interactions. These do not occur in the tree-level SM Lagrangian, but they in general may appear in the EFT with D=6 operators. We parametrize them as follows:

$$
\begin{aligned}
\mathcal{L}_{\rm dipole} = -\frac{1}{4v}\Bigg[& g_s\sum_{f\in u,d}\frac{\sqrt{m_{f_i}m_{f_j}}}{v}\bar f_i\sigma_{\mu\nu}T^a[d_{Gf}]_{ij}f_jG^a_{\mu\nu} + e\sum_{f\in u,d,e}\frac{\sqrt{m_{f_i}m_{f_j}}}{v}\bar f_i\sigma_{\mu\nu}[d_{Af}]_{ij}f_jA_{\mu\nu}\\
&+\sqrt{g^2+g'^2}\sum_{f\in u,d,e}\frac{\sqrt{m_{f_i}m_{f_j}}}{v}\bar f_i\sigma_{\mu\nu}[d_{Zf}]_{ij}f_jZ_{\mu\nu}\\
&+\sqrt{2}g\left(\frac{\sqrt{m_{u_i}m_{u_j}}}{v}\bar d_{L,i}\sigma_{\mu\nu}[d_{Wu}]_{ij}u_{R,j}W^-_{\mu\nu} + \frac{\sqrt{m_{d_i}m_{d_j}}}{v}\bar u_{L,i}\sigma_{\mu\nu}[d_{Wd}]_{ij}d_{R,j}W^+_{\mu\nu} + \text{h.c.}\right)\\
&+\sqrt{2}g\left(\frac{\sqrt{m_{e_i}m_{e_j}}}{v}\bar\nu_{L,i}\sigma_{\mu\nu}[d_{We}]_{ij}e_{R,j}W^+_{\mu\nu} + \text{h.c.}\right)\\
&+g_s\sum_{f\in u,d}\frac{\sqrt{m_{f_i}m_{f_j}}}{v}\bar f_i\sigma_{\mu\nu}T^a[\tilde d_{Gf}]_{ij}f_j\widetilde G^a_{\mu\nu} + e\sum_{f\in u,d,e}\frac{\sqrt{m_{f_i}m_{f_j}}}{v}\bar f_i\sigma_{\mu\nu}[\tilde d_{Af}]_{ij}f_j\widetilde A_{\mu\nu}\\
&+\sqrt{g^2+g'^2}\frac{\sqrt{m_{f_i}m_{f_j}}}{v}\sum_{f\in u,d,e}\bar f_i\sigma_{\mu\nu}[\tilde d_{Zf}]_{ij}f_j\widetilde Z_{\mu\nu}\Bigg],
\end{aligned}
$$

$$(1.11)$$

where $\sigma_{\mu\nu} = i[\gamma_\mu, \gamma_\nu]/2$, and d_{Af}, $\tilde{d}_{Af}$, d_{Zf}, $\tilde{d}_{Zf}$ are Hermitian 3×3 matrices, while d_{Wf} are general complex 3×3 matrices. The field strength tensors are defined as $X_{\mu\nu} = \partial_\mu X_\nu - \partial_\nu X_\mu$, and $\tilde{X}_{\mu\nu} = \epsilon_{\mu\nu\rho\sigma}\partial_\rho X_\sigma$. The coefficients d_{vf} are related to the Wilson coefficients in the Warsaw basis as

$$\begin{aligned}
d_{Gf} - i\tilde{d}_{Gf} &= -2\sqrt{2} c_{fG}, \\
d_{Af} - i\tilde{d}_{Af} &= -2\sqrt{2} \left(\eta_f c_{fW} + c_{fB}\right), \\
d_{Zf} - i\tilde{d}_{Zf} &= -\frac{2\sqrt{2}}{g^2 + g'^2} \left(g^2 \eta_f c_{fW} - g'^2 c_{fB}\right), \\
d_{Wf} &= -2\sqrt{2} c_{fW},
\end{aligned} \tag{1.12}$$

where $\eta_u = +1$, $\eta_{d,e} = -1$, and the formulas in the SILH basis are the same with $c_i \to s_i$.

Gauge boson self-interactions

Gauge boson self-interactions are not directly relevant for LHC Higgs searches, however we include them in this presentation because of the important synergy between the triple gauge couplings and Higgs couplings measurements [6,20–24]. The triple gauge interactions in the effective Lagrangian are parameterized by

$$\begin{aligned}
\mathcal{L}_{\rm tgc} &= ie \left(W^+_{\mu\nu} W^-_\mu - W^-_{\mu\nu} W^+_\mu\right) A_\nu + ie \left[(1 + \delta\kappa_\gamma) A_{\mu\nu}\, W^+_\mu W^-_\nu + \tilde{\kappa}_\gamma \tilde{A}_{\mu\nu}\, W^+_\mu W^-_\nu\right] \\
&+ igc_\theta \left[(1 + \delta g_{1,z}) \left(W^+_{\mu\nu} W^-_\mu - W^-_{\mu\nu} W^+_\mu\right) Z_\nu + (1 + \delta\kappa_z)\, Z_{\mu\nu}\, W^+_\mu W^-_\nu + \tilde{\kappa}_z \tilde{Z}_{\mu\nu}\, W^+_\mu W^-_\nu\right] \\
&+ i\frac{e}{m_W^2} \left[\lambda_\gamma W^+_{\mu\nu} W^-_{\nu\rho} A_{\rho\mu} + \tilde{\lambda}_\gamma W^+_{\mu\nu} W^-_{\nu\rho} \tilde{A}_{\rho\mu}\right] + i\frac{gc_\theta}{m_W^2} \left[\lambda_z W^+_{\mu\nu} W^-_{\nu\rho} Z_{\rho\mu} + \tilde{\lambda}_z W^+_{\mu\nu} W^-_{\nu\rho} \tilde{Z}_{\rho\mu}\right] \\
&- g_s f^{abc} \partial_\mu G^a_\nu G^b_\mu G^c_\nu. + \frac{c_{3G}}{v^2} g_s^3 f^{abc} G^a_{\mu\nu} G^b_{\nu\rho} G^c_{\rho\mu} + \frac{\tilde{c}_{3G}}{v^2} g_s^3 f^{abc} \widetilde{G}^a_{\mu\nu} G^b_{\nu\rho} G^c_{\rho\mu}.
\end{aligned} \tag{1.13}$$

The couplings of electroweak gauge bosons follow the customary parametrization of Ref. [25]. The anomalous triple gauge couplings of electroweak gauge bosons are related to the Wilson coefficients in the Warsaw basis as

$$\begin{aligned}
\delta g_{1,z} &= \frac{g^2 + g'^2}{g^2 - g'^2} \left(-g'^2 c_{WB} + c_T - \delta v\right), \\
\delta\kappa_\gamma &= g^2 c_{WB}, \\
\delta\kappa_z &= -2c_{WB} \frac{g^2 g'^2}{g^2 - g'^2} + \frac{g^2 + g'^2}{g^2 - g'^2} \left(c_T - \delta v\right), \\
\lambda_\gamma &= -\frac{3}{2} g^4 c_{3W}, \\
\lambda_z &= -\frac{3}{2} g^4 c_{3W}, \\
\tilde{\kappa}_\gamma &= g^2 \tilde{c}_{WB}, \\
\tilde{\kappa}_z &= -g'^2 \tilde{c}_{WB}, \\
\tilde{\lambda}_\gamma &= -\frac{3}{2} g^4 \tilde{c}_{3W}, \\
\tilde{\lambda}_z &= -\frac{3}{2} g^4 \tilde{c}_{3W}.
\end{aligned} \tag{1.14}$$

The analogous relations for the SILH basis read

$$\begin{aligned}
\delta g_{1z} &= -\frac{g^2 + g'^2}{4(g^2 - g'^2)} \left[(g^2 - g'^2) s_{HW} + g^2 (s_W + s_{2W}) + g'^2 (s_B + s_{2B}) - 4 s_T + 2 [s'_{H\ell}]_{22}\right], \\
\delta\kappa_\gamma &= -\frac{g^2}{4} \left[s_{HW} + s_{HB}\right], \\
\delta\kappa_z &= -\frac{1}{4} \left(g^2 s_{HW} - g'^2 s_{HB}\right) - \frac{g^2 + g'^2}{4(g^2 - g'^2)} \left[g^2 (s_W + s_{2W}) + g'^2 (s_B + s_{2B}) - 4 s_T + 2 [s'_{H\ell}]_{22}\right], \\
\lambda_z &= -\frac{3}{2} g^4 s_{3W}, \qquad \lambda_\gamma = \lambda_z, \\
\delta\tilde{\kappa}_\gamma &= -\frac{g^2}{4} \left[\tilde{s}_{HW} + \tilde{s}_{HB}\right],
\end{aligned}$$

$$
\begin{aligned}
\delta\tilde{\kappa}_z &= \frac{g'^2}{4}\left[\tilde{s}_{HW}+\tilde{s}_{HB}\right], \\
\tilde{\lambda}_z &= -\frac{3}{2}g^4\tilde{s}_{3W}, \qquad \tilde{\lambda}_\gamma = \tilde{\lambda}_z.
\end{aligned}
\tag{1.15}
$$

Single Higgs couplings

In this subsection we discuss the terms in the effective Lagrangian that involve a single Higgs boson field h. This part is the most relevant one from the point of view of the LHC Higgs phenomenology.

We first define the Higgs boson couplings to a pair of fermions:

$$
\mathcal{L}_{\rm hff} = -\frac{h}{v}\sum_{f\in u,d,e}\sum_{ij}\sqrt{m_{f_i}m_{f_j}}\left(\delta_{ij}+[\delta y_f]_{ij}e^{i[\phi_f]_{ij}}\right)\bar{f}_{R,i}f_{L,j}+\text{h.c.},
\tag{1.16}
$$

where $[\delta y_f]_{ij}$ and ϕ_{ij} are general 3×3 matrices with real element. The corrections to the SM Yukawa interactions are related to the Wilson coefficients in the Warsaw and SILH basis by

$$
\begin{aligned}
[\delta y_f]_{ij}e^{i[\phi_f]_{ij}} &= \frac{1}{\sqrt{2}}[c_f]_{ij}-\delta_{ij}\left(c_H+\delta v\right) \\
&= [s_f]_{ij}-\delta_{ij}\left[s_H+\frac{1}{2}[s'_{H\ell}]_{22}\right],
\end{aligned}
\tag{1.17}
$$

Next, we define the following single Higgs boson couplings to a pair of the SM gauge fields:

$$
\begin{aligned}
\mathcal{L}_{\rm hvv} = \frac{h}{v}\Bigg[&(1+\delta c_w)\frac{g^2v^2}{2}W^+_\mu W^-_\mu+(1+\delta c_z)\frac{(g^2+g'^2)v^2}{4}Z_\mu Z_\mu \\
- \ &+c_{ww}\frac{g^2}{2}W^+_{\mu\nu}W^-_{\mu\nu}+\tilde{c}_{ww}\frac{g^2}{2}W^+_{\mu\nu}\tilde{W}^-_{\mu\nu}+c_{w\Box}g^2\left(W^-_\mu\partial_\nu W^+_{\mu\nu}+\text{h.c.}\right) \\
&+c_{gg}\frac{g_s^2}{4}G^a_{\mu\nu}G^a_{\mu\nu}+c_{\gamma\gamma}\frac{e^2}{4}A_{\mu\nu}A_{\mu\nu}+c_{z\gamma}\frac{e\sqrt{g^2+g'^2}}{2}Z_{\mu\nu}A_{\mu\nu}+c_{zz}\frac{g^2+g'^2}{4}Z_{\mu\nu}Z_{\mu\nu} \\
&+c_{z\Box}g^2Z_\mu\partial_\nu Z_{\mu\nu}+c_{\gamma\Box}gg'Z_\mu\partial_\nu A_{\mu\nu} \\
&+\tilde{c}_{gg}\frac{g_s^2}{4}G^a_{\mu\nu}\tilde{G}^a_{\mu\nu}+\tilde{c}_{\gamma\gamma}\frac{e^2}{4}A_{\mu\nu}\tilde{A}_{\mu\nu}+\tilde{c}_{z\gamma}\frac{e\sqrt{g^2+g'^2}}{2}Z_{\mu\nu}\tilde{A}_{\mu\nu}+\tilde{c}_{zz}\frac{g^2+g'^2}{4}Z_{\mu\nu}\tilde{Z}_{\mu\nu}\Bigg],
\end{aligned}
\tag{1.18}
$$

where all the couplings above are real. The terms in the first two lines describe corrections to the SM Higgs couplings to W and Z, while the remaining terms introduce Higgs couplings to gauge bosons with a tensor structure that is absent in the SM Lagrangian. Note that, using equations of motion, we could get rid of certain 2-derivative interactions between the Higgs and gauge bosons: $hZ_\mu\partial_\nu Z_{\nu\mu}$, $hZ_\mu\partial_\nu A_{\nu\mu}$, and $hW^\pm_\mu\partial_\nu W^\mp_{\nu\mu}$. These interactions would then be traded for contact interactions of the Higgs, gauge bosons and fermions in Eq. (1.5). However, one of the defining features of our effective Lagrangian is that the coefficients of the latter couplings are equal to the corresponding vertex correction in Eq. (1.5). This form can be always obtained, without any loss of generality, starting from an arbitrary dimension-6 Lagrangian provided the 2-derivative $hV_\mu\partial_\nu V_{\nu\mu}$ are kept in the Lagrangian. Note that we work in the limit where the neutrinos are massless and the Higgs boson does not couple to the neutrinos. In the EFT context, the couplings to neutrinos induced by dimension-5 operators are proportional to neutrino masses, therefore they are far too small to have any relevance for LHC phenomenology.

The shifts of the Higgs couplings to W and Z bosons are related to the Wilson coefficients in the Warsaw and SILH basis by

$$
\begin{aligned}
\delta c_w &= -c_H-c_{WB}\frac{4g^2g'^2}{g^2-g'^2}+4c_T\frac{g^2}{g^2-g'^2}-\delta v\frac{3g^2+g'^2}{g^2-g'^2} \\
&= -s_H-\frac{g^2g'^2}{g^2-g'^2}\left[s_W+s_B+s_{2W}+s_{2B}-\frac{4}{g'^2}s_T+\frac{3g^2+g'^2}{2g^2g'^2}[s'_{H\ell}]_{22}\right], \\
\delta c_z &= -c_H-3\delta v \\
&= -s_H-\frac{3}{2}[s'_{H\ell}]_{22},
\end{aligned}
\tag{1.19}
$$

The two-derivative Higgs couplings to gauge bosons are related to the Wilson coefficients in the Warsaw basis by

$$
c_{gg} = c_{GG},
$$

$$
\begin{aligned}
c_{\gamma\gamma} &= c_{WW} + c_{BB} - 4c_{WB}, \\
c_{zz} &= \frac{g^4 c_{WW} + g'^4 c_{BB} + 4g^2 g'^2 c_{WB}}{(g^2+g'^2)^2}, \\
c_{z\Box} &= -\frac{2}{g^2}\left(c_T - \delta v\right), \\
c_{z\gamma} &= \frac{g^2 c_{WW} - g'^2 c_{BB} - 2(g^2-g'^2) c_{WB}}{g^2+g'^2}, \\
c_{\gamma\Box} &= \frac{2}{g^2-g'^2}\left((g^2+g'^2)c_{WB} - 2c_T + 2\delta v\right), \\
c_{ww} &= c_{WW}, \\
c_{w\Box} &= \frac{2}{g^2-g'^2}\left(g'^2 c_{WB} - c_T + \delta v\right).
\end{aligned}
\tag{1.20}
$$

and the same for the CP-odd couplings $\tilde{c}_{gg}$, $\tilde{c}_{\gamma\gamma}$, $\tilde{c}_{z\gamma}$, $\tilde{c}_{zz}$, $\tilde{c}_{ww}$, with $c \to \tilde{c}$ on the right hand side. The analogous expressions for the SILH basis read

$$
\begin{aligned}
c_{gg} &= s_{GG}, \\
c_{\gamma\gamma} &= s_{BB}, \\
c_{zz} &= -\frac{1}{g^2+g'^2}\left[g^2 s_{HW} + g'^2 s_{HB} - g'^2 s_\theta^2 s_{BB}\right], \\
c_{z\Box} &= \frac{1}{2g^2}\left[g^2(s_W + s_{HW} + s_{2W}) + g'^2(s_B + s_{HB} + s_{2B}) - 4s_T + 2[s'_{H\ell}]_{22}\right], \\
c_{z\gamma} &= \frac{s_{HB} - s_{HW}}{2} - s_\theta^2 s_{BB}, \\
c_{\gamma\Box} &= \frac{s_{HW} - s_{HB}}{2} + \frac{1}{g^2-g'^2}\left[g^2(s_W + s_{2W}) + g'^2(s_B + s_{2B}) - 4s_T + 2[s'_{H\ell}]_{22}\right], \\
c_{ww} &= -s_{HW}, \\
c_{w\Box} &= \frac{s_{HW}}{2} + \frac{1}{2(g^2-g'^2)}\left[g^2(s_W + s_{2W}) + g'^2(s_B + s_{2B}) - 4s_T + 2[s'_{H\ell}]_{22}\right],
\end{aligned}
\tag{1.21}
$$

Next, couplings of the Higgs boson to a gauge field and two fermions (which are not present in the SM Lagrangian) can be generated by dimension-6 operators. The vertex-like contact interactions between the Higgs, electroweak gauge bosons, and fermions are parametrized as:

$$
\begin{aligned}
\mathcal{L}_{hvff} &= \sqrt{2} g \frac{h}{v} W_\mu^+ \left(\bar{u}_L \gamma_\mu \delta g_L^{hWq} d_L + \bar{u}_R \gamma_\mu \delta g_R^{hWq} d_R + \bar{\nu}_L \gamma_\mu \delta g_L^{hW\ell} e_L\right) + \text{h.c.} \\
&+ 2\frac{h}{v}\sqrt{g^2+g'^2} Z_\mu \left[\sum_{f=u,d,e,\nu} \bar{f}_L \gamma_\mu \delta g_L^{hZf} f_L + \sum_{f=u,d,e} \bar{f}_R \gamma_\mu \delta g_R^{hZf} f_R\right],
\end{aligned}
\tag{1.22}
$$

By construction (condition #5), the coefficients of these interaction are equal to the corresponding vertex correction in Eq. (1.5):

$$
\delta g^{hzf} = \delta g^{Zf}, \qquad \delta g^{hWf} = \delta g^{Wf}. \tag{1.23}
$$

The dipole-type contact interactions of the Higgs boson are parametrized as:

$$
\begin{aligned}
\mathcal{L}_{\text{hdvff}} = -\frac{h}{4v^2}\Bigg[& g_s \sum_{f\in u,d} \bar{f}\sigma_{\mu\nu} T^a d_{hGf} f G_{\mu\nu}^a + e \sum_{f\in u,d,e} \bar{f}\sigma_{\mu\nu} d_{hAf} f A_{\mu\nu} + \sqrt{g^2+g'^2} \sum_{f\in u,d,e} \bar{f}\sigma_{\mu\nu} d_{hZf} f Z_{\mu\nu} \\
& + \sqrt{2} g \left(\bar{d}_L \sigma_{\mu\nu} d_{hWu} u_R W_{\mu\nu}^- + \bar{u}_L \sigma_{\mu\nu} d_{hWd} d_R W_{\mu\nu}^+ + \bar{\nu}_L \sigma_{\mu\nu} d_{hWe} e_R W_{\mu\nu}^+ + \text{h.c.}\right) \\
& + g_s \sum_{f\in u,d} \bar{f}\sigma_{\mu\nu} T^a \tilde{d}_{hGf} f \widetilde{G}_{\mu\nu}^a + e \sum_{f\in u,d,e} \bar{f}\sigma_{\mu\nu} \tilde{d}_{hAf} f \widetilde{A}_{\mu\nu} + \sqrt{g^2+g'^2} \sum_{f\in u,d,e} \bar{f}\sigma_{\mu\nu} \tilde{d}_{hZf} f \widetilde{Z}_{\mu\nu}\Bigg],
\end{aligned}
\tag{1.24}
$$

where d_{hAf}, $\tilde{d}_{hAf}$, d_{hZf}, $\tilde{d}_{hZf}$ are Hermitian 3×3 matrices, while d_{hWf} are general complex 3×3 matrices. The coefficients are simply related to the corresponding dipole interactions in Eq. (1.11):

$$
d_{hVf} = d_{Vf}. \tag{1.25}
$$

Dimension-6 operators can also induce single Higgs couplings to 3 gauge bosons, but we do not display them here.

Higgs boson self-couplings and double Higgs couplings

The cubic Higgs boson self-coupling and couplings of two Higgs boson fields to matter play a role in the EFT description of double Higgs production [26, 27]. The cubic Higgs boson self-coupling is parametrized as

$$\mathcal{L}_{h,\mathrm{self}} = -(\lambda + \delta\lambda_3) v h^3. \tag{1.26}$$

The relation between the cubic Higgs coupling correction and the Wilson coefficients in the Warsaw and SILH basis is given by

$$\begin{aligned} \delta\lambda_3 &= -\lambda\,(3c_H + \delta v) - c_{6H} \\ &= -\lambda\left(3s_H + \frac{1}{2}[s'_{H\ell}]_{22}\right) - s_{6H}. \end{aligned} \tag{1.27}$$

In accordance with the condition #4, the 2-derivative Higgs boson self-couplings have been traded for other equivalent interactions and do not occur in the mass eigenstate Lagrangian. Self-interactions terms with 4, 5, and 6 Higgs boson fields may also arise from dimension-6 operators, but we do not display them here.

The interactions between two Higgs bosons and two other SM fields are parametrized as follows:

$$\begin{aligned} \mathcal{L}_{hh} &= h^2\left(1 + 2\delta c_z^{(2)}\right)\frac{g^2 + g'^2}{4} Z_\mu Z_\mu + h^2\left(1 + 2\delta c_w^{(2)}\right)\frac{g^2}{2} W_\mu^+ W_\mu^- - \frac{h^2}{2v^2}\sum_{f;ij}\sqrt{m_{f_i} m_{f_j}}\left[\bar{f}_{i,R}[y_f^{(2)}]_{ij} f_{j,L} + \mathrm{h.c.}\right]. \\ &+ \frac{h^2}{8v^2}\left(c_{gg}^{(2)} g_s^2 G_{\mu\nu}^a G_{\mu\nu}^a + 2c_{ww}^{(2)} g^2 W_{\mu\nu}^+ W_{\mu\nu}^- + c_{zz}^{(2)}(g^2 + g'^2) Z_{\mu\nu} Z_{\mu\nu} + 2c_{z\gamma}^{(2)} g g' Z_{\mu\nu} A_{\mu\nu} + c_{\gamma\gamma}^{(2)} e^2 A_{\mu\nu} A_{\mu\nu}\right) \\ &+ \frac{h^2}{8v^2}\left(\tilde{c}_{gg}^{(2)} g_s^2 G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a + 2\tilde{c}_{ww}^{(2)} g^2 W_{\mu\nu}^+ \tilde{W}_{\mu\nu}^- + \tilde{c}_{zz}^{(2)}(g^2 + g'^2) Z_{\mu\nu} \tilde{Z}_{\mu\nu} + 2\tilde{c}_{z\gamma}^{(2)} g g' Z_{\mu\nu} \tilde{A}_{\mu\nu} + \tilde{c}_{\gamma\gamma}^{(2)} e^2 A_{\mu\nu} \tilde{A}_{\mu\nu}\right) \\ &- \frac{h^2}{2v^2}\left(g^2 c_{w\Box}^{(2)}(W_\mu^+ \partial_\nu W_{\nu\mu}^- + W_\mu^- \partial_\nu W_{\nu\mu}^+) + g^2 c_{z\Box}^{(2)} Z_\mu \partial_\nu Z_{\nu\mu} + g g' c_{\gamma\Box}^{(2)} Z_\mu \partial_\nu A_{\nu\mu}\right). \end{aligned} \tag{1.28}$$

All double Higgs couplings arising from D=6 operators can be expressed by the single Higgs couplings:

$$\begin{aligned} \delta c_z^{(2)} &= \delta c_z, \qquad \delta c_w^{(2)} = \delta c_z + 3\delta m, \\ [y_f^{(2)}]_{ij} &= 3[\delta y_f]_{ij} e^{i\phi_{ij}} - \delta c_z\,\delta_{ij}, \\ c_{vv}^{(2)} &= c_{vv}, \qquad \tilde{c}_{vv}^{(2)} = \tilde{c}_{vv}, \qquad v \in \{g, w, z, \gamma\}, \\ c_{v\Box}^{(2)} &= c_{v\Box}, \qquad v \in \{w, z, \gamma\}. \end{aligned} \tag{1.29}$$

Other interaction terms with two Higgs bosons involve at least 5 fields: e.g the h^2V^3 or $h^2 ffV$ contact interactions, and are not displayed here.

Other terms

In this section we wrote down the interaction terms of mass eigenstates in the D=6 EFT Lagrangian which are most relevant for LHC Higgs phenomenology. They either enter the single and double Higgs production at tree level, or they affect electroweak precision observables that are complementary to Higgs couplings measurements. The remaining terms in the mass eigenstate Lagrangian, which are not explicitly displayed in this chapter, are contained in $\mathcal{L}_{\mathrm{other}}$ in Eq. (1.1). They include 4-fermion terms, couplings of a single Higgs boson to 3 or more gauge bosons, quartic Higgs and gauge boson self-interactions, dipole-like interactions of two gauge bosons and two fermions, and interaction terms with 5 or more fields.

3 Higgs basis

In the previous section we related the Wilson coefficients in the Warsaw and SILH bases of D=6 operators to the couplings of mass eigenstates in the Lagrangian. With this information at hand, one can proceed to calculating observables at a given order in the EFT as a function of the Wilson coefficients. The information provided above is enough to calculate the leading order EFT corrections to SM predictions for single and double Higgs production and decays in all phenomenologically relevant channels.

There is no theoretical obstacle to present the results of LHC Higgs analyses as constraints on the Wilson coefficients in the Warsaw or SILH basis. However, this procedure may not be the most efficient one from the experimental point of view. The reason is that the relation between these Wilson coefficients and the couplings of the Higgs boson in the Lagrangian is somewhat complicated, c.f. Eqs (1.6), (1.19), (1.20). In this section we propose another, equivalent parametrization of the EFT with D=6 operators. The idea, put forward in Ref. [28], is to parametrize the space of D=6 operators using a subset of couplings in the mass eigenstate Lagrangian, such as the one defined in Eq. (1.1) of Section 2. The parametrization described in this section, which differs slightly from that in Ref. [28], is referred to as the *Higgs basis*.

The salient features of the Higgs basis are the following. The goal is to parametrize the space of D=6 operators in a way that can be more directly connected to observable quantities in Higgs physics. Technically, the Higgs basis can be defined as a linear transformation from the Warsaw or SILH basis into the coefficients of certain interaction terms of the mass eigenstates (in particular the W, Z, and the Higgs bosons) in the effective Lagrangian. In practice, we will define the Higgs basis by choosing a subset of the couplings parametrizing interaction terms in the mass eigenstate Lagrangian in Eq. (1.1). All couplings in the subset have to be independent, in the sense that none can be expressed by the remaining ones at the level of a general $D = 6$ EFT Lagrangian. It is also a maximal such subset, which implies that their number is the same as the number of independent operators in the Warsaw or SILH basis. We will refer to this set as the *independent couplings*. They parametrize all possible deformations of the SM Lagrangian in the presence of D=6 operators. Therefore, they can be used on par with any other basis to describe the effects of dimension-6 operators on any physical observables (also those unrelated to Higgs physics). By definition of the Higgs basis, the independent couplings include single Higgs boson couplings to gauge bosons and fermions. Thanks to that, the parameters of the Higgs basis can be connected in a more intuitive way to LHC Higgs observables calculated at leading order in the EFT. Furthermore, the vertex corrections to the Z boson interactions with fermions are chosen to be among the independent couplings. As a consequence, combining experimental information from Higgs and electroweak precision observables is more transparent in the Higgs basis.

3.1 Independent couplings

We now describe the choice of independent couplings which defines the Higgs basis.

The first group of independent couplings that we parametrizes the interactions of the Higgs boson with the SM gauge boson, fermions, and with itself:

$$\begin{gathered} c_{gg},\ \delta c_z,\ c_{\gamma\gamma},\ c_{z\gamma},\ c_{zz},\ c_{z\Box},\ \tilde{c}_{gg},\ \tilde{c}_{\gamma\gamma},\ \tilde{c}_{z\gamma},\ \tilde{c}_{zz}, \\ \delta y_u,\ \delta y_d,\ \delta y_e,\ \phi_u,\ \phi_d,\ \phi_\ell,\ \delta\lambda_3. \end{gathered} \tag{1.1}$$

The first line is defined by Eq. (1.18), and the second one by Eq. (1.16) except for the last coupling which is defined in Eq. (1.28). All these couplings affect the Higgs boson production and/or decay at the leading order in the EFT. Therefore they are of crucial importance for LHC Higgs phenomenology. Moreover, at the leading order, they are not constrained at all by LEP-1 electroweak precision tests or low-energy precision observables.

The second group of independent couplings parametrizes the W boson mass and the Z and W boson couplings to fermions:

$$\begin{gathered} \delta m,\ \delta g_L^{Ze},\ \delta g_R^{Ze},\ \delta g_L^{W\ell},\ \delta g_L^{Zu},\ \delta g_R^{Zu},\ \delta g_L^{Zd},\ \delta g_R^{Zd},\ \delta g_R^{Wq}, \\ d_{Gu},\ d_{Gd},\ d_{Ae},\ d_{Au},\ d_{Ad},\ d_{Ze},\ d_{Zu},\ d_{Zd},\ \tilde{d}_{Gu},\ \tilde{d}_{Gd},\ \tilde{d}_{Ae},\ \tilde{d}_{Au},\ \tilde{d}_{Ad},\ \tilde{d}_{Ze},\ \tilde{d}_{Zu},\ \tilde{d}_{Zd}. \end{gathered} \tag{1.2}$$

Here the mass correction δm is defined in Eq. (1.2), the vertex corrections δg^i are defined in Eq. (1.5), and the dipole moments d_i are defined in Eq. (1.11). All these parameters also affect the Higgs boson production and/or decay at the leading order in the EFT. However, as opposed to the ones in Eq. (1.1), they affect at the same order electroweak and/or low-energy precision observables.

The third group of independent couplings parametrizes the self-couplings of gauge bosons:

$$\lambda_z,\ \tilde{\lambda}_z,\ c_{3G},\ \tilde{c}_{3G}. \tag{1.3}$$

They are defined in Eq. (1.13). These couplings do not affect Higgs production and decay at the leading order in EFT.

To complete the definition of the Higgs basis, one has to select the independent couplings corresponding to 4-fermion operators. We choose to parametrize them by the same set of Wilson coefficients as in the Warsaw basis:

$$\begin{gathered} c_{\ell\ell},\ c_{qq},\ c'_{qq},\ c_{\ell q},\ c'_{\ell q},\ c_{quqd},\ c'_{quqd},\ c_{\ell equ},\ c'_{\ell equ},\ c_{\ell edq}, \\ c_{\ell e},\ c_{\ell u},\ c_{\ell d},\ c_{qe},\ c_{qu},\ c'_{qu},\ c_{qd},\ c'_{qd},\ c_{ee},\ c_{uu},\ c_{dd},\ c_{eu},\ c_{ed},\ c_{ud},\ c'_{ud}. \end{gathered} \tag{1.4}$$

The parameters c_{ff} have 4 flavor indices. The non-trivial question of which combination of flavor indices constitutes an independent set was worked out in Ref. [29]. In the Higgs basis we take the same choice of independent 4-fermion couplings as in that reference, with one exception. As explained in the next subsection, in a D=6 EFT Lagrangian, the coupling $[c_{\ell\ell}]_{1221}$ multiplying a particular 4-lepton operator can be expressed by δm and δg^i. Therefore $[c_{\ell\ell}]_{1221}$ is not among the independent couplings defining the Higgs basis.

3.2 Dependent couplings

The number of parameters characterizing departure from the SM Lagrangian in Eq. (1.1) is larger than the number of Wilson coefficients in a basis of D=6 operators. Due to this fact, there must be relations among these parameters. Working in the Higgs basis, some of the parameters in the mass eigenstate Lagrangian can be expressed by the independent couplings; we call them the *dependent* couplings. The relations between dependent and independent couplings can be inferred from the matching between the effective Lagrangian and the Warsaw or SILH basis in Section 2. These relations *hold at the level of the dimension-6 Lagrangian*, and they are in general not respected in the presence of dimension-8 and higher operators.

We start with the dependent couplings in Eq. (1.18) parametrizing the single Higgs boson interactions with gauge bosons. They can be expressed in terms of the independent couplings as[4]

$$\begin{aligned}
\delta c_w &= \delta c_z + 4\delta m, \\
c_{ww} &= c_{zz} + 2 s_\theta^2 c_{z\gamma} + s_\theta^4 c_{\gamma\gamma}, \\
\tilde{c}_{ww} &= \tilde{c}_{zz} + 2 s_\theta^2 \tilde{c}_{z\gamma} + s_\theta^4 \tilde{c}_{\gamma\gamma}, \\
c_{w\Box} &= \frac{1}{g^2 - g'^2}\left[g^2 c_{z\Box} + g'^2 c_{zz} - e^2 s_\theta^2 c_{\gamma\gamma} - (g^2 - g'^2) s_\theta^2 c_{z\gamma}\right], \\
c_{\gamma\Box} &= \frac{1}{g^2 - g'^2}\left[2 g^2 c_{z\Box} + (g^2 + g'^2) c_{zz} - e^2 c_{\gamma\gamma} - (g^2 - g'^2) c_{z\gamma}\right].
\end{aligned} \tag{1.5}$$

The coefficients of W-boson dipole interactions in Eq. (1.11) are related to those of the Z and the photon as

$$\eta_f d_{wf} = d_{zf} - i\tilde{d}_{zf} + s_\theta^2 (d_{Af} - i\tilde{d}_{Af}), \tag{1.6}$$

where $\eta_u = 1$ and $\eta_{d,e} = -1$. The coefficients of the dipole-like Higgs couplings in Eq. (1.24) are simply related to the corresponding dipole moments:

$$d_{hvf} = d_{vf}, \quad \tilde{d}_{hvf} = \tilde{d}_{vf}, \qquad v \in \{g, w, z, \gamma\}. \tag{1.7}$$

Coefficients of all interaction terms with two Higgs bosons in Eq. (1.28) are dependent couplings. The can be expressed in terms of the independent couplings as:

$$\begin{aligned}
\delta c_z^{(2)} &= \delta c_z, \qquad \delta c_w^{(2)} = \delta c_z + 3\delta m, \\
[y_f^{(2)}]_{ij} &= 3[\delta y_f]_{ij} e^{i\phi_{ij}} - \delta c_z\, \delta_{ij}, \\
c_{vv}^{(2)} &= c_{vv}, \qquad \tilde{c}_{vv}^{(2)} = \tilde{c}_{vv}, \qquad v \in \{g, w, z, \gamma\}, \\
c_{v\Box}^{(2)} &= c_{v\Box}, \qquad v \in \{w, z, \gamma\}.
\end{aligned} \tag{1.8}$$

The dependent vertex corrections are expressed in terms of the independent couplings as

$$\delta g_L^{Z\nu} = \delta g_L^{Ze} + \delta g_L^{W\ell}, \qquad \delta g_L^{Wq} = \delta g_L^{Zu} V_{\rm CKM} - V_{\rm CKM} \delta g_L^{Zd}. \tag{1.9}$$

All but two triple gauge couplings in Eq. (1.13) are dependent couplings expressed in terms of the independent couplings as

$$\begin{aligned}
\delta g_{1,z} &= \frac{1}{2(g^2 - g'^2)}\left[c_{\gamma\gamma} e^2 g'^2 + c_{z\gamma}(g^2 - g'^2) g'^2 - c_{zz}(g^2 + g'^2) g'^2 - c_{z\Box}(g^2 + g'^2) g^2\right] \\
\delta\kappa_\gamma &= -\frac{g^2}{2}\left(c_{\gamma\gamma}\frac{e^2}{g^2 + g'^2} + c_{z\gamma}\frac{g^2 - g'^2}{g^2 + g'^2} - c_{zz}\right), \\
\tilde{\kappa}_\gamma &= -\frac{g^2}{2}\left(\tilde{c}_{\gamma\gamma}\frac{e^2}{g^2 + g'^2} + \tilde{c}_{z\gamma}\frac{g^2 - g'^2}{g^2 + g'^2} - \tilde{c}_{zz}\right), \\
\delta\kappa_z &= \delta g_{1,z} - t_\theta^2 \delta\kappa_\gamma, \qquad \tilde{\kappa}_z = -t_\theta^2 \tilde{\kappa}_\gamma, \\
\lambda_\gamma &= \lambda_z, \quad \tilde{\lambda}_\gamma = \tilde{\lambda}_z.
\end{aligned} \tag{1.10}$$

Finally, we discuss how the Wilson coefficient $[c_{\ell\ell}]_{1221}$ is expressed by the independent couplings. One defining feature of the mass eigenstate Lagrangian Eq. (1.1) is that the tree-level relations between the SM electroweak parameters

[4] The relation between c_{ww}, $\tilde{c}_{ww}$ and other parameters can also be viewed as a consequence of the accidental custodial symmetry at the level of the dimension-6 operators [19].

and input observables are not affected by D=6 operators (condition # 2). On the other hand, one of the four-fermion couplings in the Lagrangian,

$$\mathcal{L}_{4f}^{D=6} \supset [c_{\ell\ell}]_{1221}(\bar{\ell}_{1,L}\gamma_\rho \ell_{2,L})(\bar{\ell}_{2,L}\gamma_\rho \ell_{1,L}) \tag{1.11}$$

does affect the relation between the parameter v and the muon decay width from which $v = (\sqrt{2}G_F)^{-2}$ is determined:

$$\frac{\Gamma(\mu \to e\nu\nu)}{\Gamma(\mu \to e\nu\nu)_{\rm SM}} \approx 1 + 2[\delta g_L^{We}]_{11} + 2[\delta g_L^{We}]_{22} - 4\delta m - [c_{\ell\ell}]_{1221}. \tag{1.12}$$

Therefore, the muon decay width is unchanged with respect to the SM when $[c_{\ell\ell}]_{1221}$ is related to δm and δg as

$$[c_{\ell\ell}]_{1221} = 2\delta[g_L^{We}]_{11} + 2[\delta g_L^{We}]_{22} - 4\delta m. \tag{1.13}$$

In other words, due to the fact that we selected δm and δg selected as an independent coupling in the Higgs basis, $[c_{\ell\ell}]_{1221}$ has to be a dependent coupling. Of course, one could equivalently choose $[c_{\ell\ell}]_{1221}$ to define a basis, and remove e.g. δm from the list of independent couplings.

3.3 Summary and comments

In summary, in the Higgs basis the parameters spanning the space of D=6 EFT operators are the independent couplings in Eqs. (1.1), (1.2), (1.3), and (1.4). In the EFT expansion, the independent couplings are formally of order $\mathcal{O}(\Lambda^{-2}$. These parameters describe certain deviations from the SM interactions in the mass eigenstate Lagrangian in Eq. (1.1). All other deviations in the mass eigenstate Lagrangian can be expressed by the independent coupling.

The Higgs basis can be used in par with any other basis to describe the effects of dimension-6 operators on physical observables. It should be stressed that it is not intrinsically better or worse than any other complete basis. Its usefulness is in the fact that description of Higgs observables and electroweak precision observables at the leading EFT order (tree-level $\mathcal{O}(\Lambda^{-2})$) is more transparent than in other bases. On the other hand, most of the existing one-loop EFT calculations have been performed in the SILH [30–33] Warsaw [29, 34–38] basis, therefore these bases are currently the natural choice as far as analyses beyond the leading order are concerned. Nevertheless, experimental constraints on the parameters in the Higgs basis can be always translated to other bases. To this end, the linear map between the parameters in the Higgs basis and the Wilson coefficient in the SILH and Warsaw bases provided in Section 2 can used (see e.g. [24] for the translation of the LHC Higgs and TGC constraints). These maps are used by the Rosetta program [39], which provides automated translation between different bases and an interface to Monte Carlo simulations in the MadGraph 5 framework [40]. At the same time, the independent couplings can be easily connected to Higgs *pseudo-observables* at the amplitude level, as defined e.g. in Ref. [41].

In total, the Higgs basis, as any complete basis at the dimension-6 level, is parametrized by 2499 independent real couplings [29]. One should not, however, be intimidated by this number. The point is that a much smaller subset of the independent couplings is relevant for analyses of Higgs data at leading EFT order. First of all, the coefficients of 4-fermion interactions in Eq. (1.4) and triple gauge interactions in Eq. (1.3) do not enter Higgs observables at the leading order. At that order, the parameters relevant for LHC Higgs analyses are those in Eqs. (1.1) and (1.2), which already reduces the number of variables by a significant number. Furthermore, there are several motivated assumptions about the UV theory underlying the EFT which could be used to further reduce the number of parameters:

– *Minimal flavor violation*, in which case the matrices δy_f, d^f, δg^f, and $\sin\phi_f$ reduce to a single number for each f.
– *CP conservation*, in which case all CP-odd couplings vanish: $\tilde{c}_i = \phi_f = \tilde{d}_f = 0$.
– *Custodial symmetry*, in which case $\delta m = 0$.[5]

We stress that independent couplings should not be arbitrarily set to zero without an underlying symmetry assumption. Furthermore, the relations between the dependent and independent couplings in the mass eigenstate Lagrangian should be consistently imposed, so as to preserve the structure of the D=6 EFT Lagrangian.

Finally, to reduce the number of variables, one can take advantage of the fact that, in addition to Higgs observables, other measurements are sensitive to the parameters in Eq. (1.2). In particular, the parameters in the first line of Eq. (1.2) are constrained by electroweak precision tests in LEP-1. These are among the most stringent constraints on EFT parameters, and they have an important impact on possible signals in Higgs searches. Assuming minimal flavor violation, all the vertex corrections in Eq. (1.2) are constrained to be smaller than $O(10^{-3})$ (for the leptonic vertex corrections and δm), or $O(10^{-2})$ (for the quark vertex corrections) [21, 42, 43].[6] Even when the assumption of minimal flavor violation is

[5] Custodial symmetry implies several relations between Higgs couplings to gauge bosons: $\delta c_w = \delta c_z$, $c_{w\Box} = c_\theta^2 c_{z\Box} + s_\theta^2 c_{\gamma\Box}$, $c_{ww} = c_{zz} + 2s_\theta^2 c_{z\gamma} + s_\theta^4 c_\gamma$, and $\tilde{c}_{ww} = \tilde{c}_{zz} + 2s_\theta^2 \tilde{c}_{z\gamma} + s_\theta^4 \tilde{c}_\gamma$. The last three are satisfied automatically at the level of dimension-6 Lagrangian, while the first one is true for $\delta m = 0$, see Eq. (1.5).

[6] These constraints may be relaxed if the D=6 EFT does not provide a good description of electroweak precision observables [44]. Such cases are discussed in more detail in Section 4.

not imposed, all the leptonic, bottom and charm quark vertex corrections are still constrained at the level of $O(10^{-2})$ or better [45]. Similarly, many parameters in the second line of Eq. (1.2) are strongly constrained by measurements of the magnetic and electric dipole moments. In the LHC environment, experimental sensitivity is often not sufficient to probe these parameters with a comparable accuracy. If that is indeed the case, it is well-motivated to neglect the parameters in Eq. (1.2) in LHC Higgs analyses.

Once the parameters in Eq. (1.2) are neglected, this leaves the parameter in Eq. (1.1) to describe Higgs observables. This set consists of 10 bosonic and $2 \times 3 \times 3 \times 3 = 54$ fermionic couplings. Furthermore, 31 of these couplings are CP-odd, therefore they affect the Higgs signal strength measurements only at the quadratic level ($\mathcal{O}(\Lambda^{-4})$ in the EFT expansion), while flavor off-diagonal Yukawa couplings only affect exotic Higgs decays. In the limit where fermionic couplings respect the minimal flavor violation paradigm, 9 parameters remain to describe leading order EFT corrections to the existing Higgs signal strength measurements at the LHC. In the Higgs basis, these 9 parameters are:

$$c_{gg},\ \delta c_z,\ c_{\gamma\gamma},\ c_{z\gamma},\ c_{zz},\ c_{z\Box},\ \delta y_u,\ \delta y_d,\ \delta y_e. \tag{1.14}$$

4 Comments of EFT validity

	Bosonic
O_H	$\left[\partial_\mu(H^\dagger H)\right]^2$
O_T	$\left(H^\dagger \overleftrightarrow{D}_\mu H\right)^2$
O_{6H}	$(H^\dagger H)^3$
O_{3G}	$g_s^3 f^{abc} G^a_{\mu\nu} G^b_{\nu\rho} G^c_{\rho\mu}$
$O_{3\widetilde{G}}$	$g_s^3 f^{abc} \widetilde{G}^a_{\mu\nu} G^b_{\nu\rho} G^c_{\rho\mu}$
O_{3W}	$g^3 \epsilon^{ijk} W^i_{\mu\nu} W^j_{\nu\rho} W^k_{\rho\mu}$
$O_{3\widetilde{W}}$	$g^3 \epsilon^{ijk} \widetilde{W}^i_{\mu\nu} W^j_{\nu\rho} W^k_{\rho\mu}$
O_{GG}	$\frac{g_s^2}{4} H^\dagger H\, G^a_{\mu\nu} G^a_{\mu\nu}$
$O_{\widetilde{GG}}$	$\frac{g_s^2}{4} H^\dagger H\, \widetilde{G}^a_{\mu\nu} G^a_{\mu\nu}$
O_{WW}	$\frac{g^2}{4} H^\dagger H\, W^i_{\mu\nu} W^i_{\mu\nu}$
$O_{\widetilde{WW}}$	$\frac{g^2}{4} H^\dagger H\, \widetilde{W}^i_{\mu\nu} W^i_{\mu\nu}$
O_{BB}	$\frac{g'^2}{4} H^\dagger H\, B_{\mu\nu} B_{\mu\nu}$
$O_{\widetilde{BB}}$	$\frac{g'^2}{4} H^\dagger H\, \widetilde{B}_{\mu\nu} B_{\mu\nu}$
O_{WB}	$g g' H^\dagger \sigma^i H\, W^i_{\mu\nu} B_{\mu\nu}$
$O_{\widetilde{WB}}$	$g g' H^\dagger \sigma^i H\, \widetilde{W}^i_{\mu\nu} B_{\mu\nu}$

Table 1.1: Bosonic dimension-6 operators in the Warsaw basis.

Yukawa and Vertex		Dipole	
$[O_e]_{ij}$	$-\frac{\sqrt{m_{e_i} m_{e_j}}}{v}(H^\dagger H - \frac{v^2}{2})\bar{e}_i H^\dagger \ell_j$	$[O_{eW}]_{ij}$	$g\frac{\sqrt{m_{e_i} m_{e_j}}}{v}\bar{\ell}_i \sigma^k H \sigma_{\mu\nu} e_j W^k_{\mu\nu}$
$[O_u]_{ij}$	$-\frac{\sqrt{m_{u_i} m_{u_j}}}{v}(H^\dagger H - \frac{v^2}{2})\bar{u}_i \widetilde{H}^\dagger q_j$	$[O_{eB}]_{ij}$	$g'\frac{\sqrt{m_{e_i} m_{e_j}}}{v}\bar{\ell}_i H \sigma_{\mu\nu} e_j B_{\mu\nu}$
$[O_d]_{ij}$	$-\frac{\sqrt{m_{d_i} m_{d_j}}}{v}(H^\dagger H - \frac{v^2}{2})\bar{d}_i H^\dagger q_j$	$[O_{uG}]_{ij}$	$g_s\frac{\sqrt{m_{u_i} m_{u_j}}}{v}\bar{q}_i \tilde{H} \sigma_{\mu\nu} T^a u_j G^a_{\mu\nu}$
$[O_{H\ell}]_{ij}$	$i\bar{\ell}_i \gamma_\mu \ell_j H^\dagger \overleftrightarrow{D}_\mu H$	$[O_{uW}]_{ij}$	$g\frac{\sqrt{m_{u_i} m_{u_j}}}{v}\bar{q}_i \sigma^k \tilde{H} \sigma_{\mu\nu} u_j W^k_{\mu\nu}$
$[O'_{H\ell}]_{ij}$	$i\bar{\ell}_i \sigma^k \gamma_\mu \ell_j H^\dagger \sigma^k \overleftrightarrow{D}_\mu H$	$[O_{uB}]_{ij}$	$g'\frac{\sqrt{m_{u_i} m_{u_j}}}{v}\bar{q}_i \tilde{H} \sigma_{\mu\nu} u_j B_{\mu\nu}$
$[O_{He}]_{ij}$	$i\bar{e}_i \gamma_\mu \bar{e}_j H^\dagger \overleftrightarrow{D}_\mu H$	$[O_{dG}]_{ij}$	$g_s\frac{\sqrt{m_{d_i} m_{d_j}}}{v}\bar{q}_i H \sigma_{\mu\nu} T^a d_j G^a_{\mu\nu}$
$[O_{Hq}]_{ij}$	$i\bar{q}_i \gamma_\mu q_j H^\dagger \overleftrightarrow{D}_\mu H$	$[O_{dW}]_{ij}$	$g\frac{\sqrt{m_{d_i} m_{d_j}}}{v}\bar{q}_i \sigma^k H \sigma_{\mu\nu} d_j W^k_{\mu\nu}$
$[O'_{Hq}]_{ij}$	$i\bar{q}_i \sigma^k \gamma_\mu q_j H^\dagger \sigma^k \overleftrightarrow{D}_\mu H$	$[O_{dB}]_{ij}$	$g'\frac{\sqrt{m_{d_i} m_{d_j}}}{v}\bar{q}_i H \sigma_{\mu\nu} d_j B_{\mu\nu}$
$[O_{Hu}]_{ij}$	$i\bar{u}_i \gamma_\mu u_j H^\dagger \overleftrightarrow{D}_\mu H$		
$[O_{Hd}]_{ij}$	$i\bar{d}_i \gamma_\mu d_j H^\dagger \overleftrightarrow{D}_\mu H$		
$[O_{Hud}]_{ij}$	$i\bar{u}_i \gamma_\mu d_j \tilde{H}^\dagger D_\mu H$		

Table 1.2: Two-fermion dimension-6 operators in the Warsaw basis. In this table, e, u, d are always right-handed fermions, while ℓ and q are left-handed. For complex operators the complex conjugate operator is implicit.

$(\bar{L}L)(\bar{L}L)$ and $(\bar{L}R)(\bar{L}R)$		$(\bar{R}R)(\bar{R}R)$		$(\bar{L}L)(\bar{R}R)$	
$O_{\ell\ell}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{\ell}\gamma_\mu \ell)$	O_{ee}	$(\bar{e}\gamma_\mu e)(\bar{e}\gamma_\mu e)$	$O_{\ell e}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{e}\gamma_\mu e)$
O_{qq}	$(\bar{q}\gamma_\mu q)(\bar{q}\gamma_\mu q)$	O_{uu}	$(\bar{u}\gamma_\mu u)(\bar{u}\gamma_\mu u)$	$O_{\ell u}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{u}\gamma_\mu u)$
O'_{qq}	$(\bar{q}\gamma_\mu \sigma^i q)(\bar{q}\gamma_\mu \sigma^i q)$	O_{dd}	$(\bar{d}\gamma_\mu d)(\bar{d}\gamma_\mu d)$	$O_{\ell d}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{d}\gamma_\mu d)$
$O_{\ell q}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{q}\gamma_\mu q)$	O_{eu}	$(\bar{e}\gamma_\mu e)(\bar{u}\gamma_\mu u)$	O_{eq}	$(\bar{q}\gamma_\mu q)(\bar{e}\gamma_\mu e)$
$O'_{\ell q}$	$(\bar{\ell}\gamma_\mu \sigma^i \ell)(\bar{q}\gamma_\mu \sigma^i q)$	O_{ed}	$(\bar{e}\gamma_\mu e)(\bar{d}\gamma_\mu d)$	O_{qu}	$(\bar{q}\gamma_\mu q)(\bar{u}\gamma_\mu u)$
O_{quqd}	$(\bar{q}^j u)\epsilon_{jk}(\bar{q}^k d)$	O_{ud}	$(\bar{u}\gamma_\mu u)(\bar{d}\gamma_\mu d)$	O'_{qu}	$(\bar{q}\gamma_\mu T^a q)(\bar{u}\gamma_\mu T^a u)$
O'_{quqd}	$(\bar{q}^j T^a u)\epsilon_{jk}(\bar{q}^k T^a d)$	O'_{ud}	$(\bar{u}\gamma_\mu T^a u)(\bar{d}\gamma_\mu T^a d)$	O_{qd}	$(\bar{q}\gamma_\mu q)(\bar{d}\gamma_\mu d)$
$O_{\ell equ}$	$(\bar{\ell}^j e)\epsilon_{jk}(\bar{q}^k u)$			O'_{qd}	$(\bar{q}\gamma_\mu T^a q)(\bar{d}\gamma_\mu T^a d)$
$O'_{\ell equ}$	$(\bar{\ell}^j \sigma_{\mu\nu} e)\epsilon_{jk}(\bar{q}^k \sigma^{\mu\nu} u)$				
$O_{\ell edq}$	$(\bar{\ell}^j e)(\bar{d} q^j)$				

Table 1.3: Four-fermion operators in the Warsaw basis [11]. In this table, e, u, d are always right-handed fermions, while ℓ and q are left-handed. A flavor index is implicit for each fermion field. For complex operators the complex conjugate operator is implicit.

Chapter 2

EFT application

1 LO EFT tools

1.1 Tools for translations (Rosetta)

1.2 Tools for calculating observables (e.g EHdecay)

1.3 Tools for simulating events (e.g. Madgraph)

1.4 Tools for comparing with experiments (e.g. Sfitter)

2 NLO EFT results

2.1 NLO EW

comparison to LO

2.2 NLO QCD

comparison to LO

3 Interpretations in terms of non-linear EFT

Part IV

Measurements and Observables[1]

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