LHCHXSWG-INT-2015-001
January 22, 2015

LHC Higgs Cross Section Working Group 2 (Higgs Properties)

Higgs Basis: Proposal for an EFT basis choice for LHC HXSWG

We acknowledge contributions and feedback from: André David, Michael Dührssen, Adam Falkowski, Christophe Grojean, Florian Görtz, Gino Isidori, Tilman Plehn, Alex Pomarol, Francesco Riva, Veronica Sanz.

This document is currently a work in progress. Acknowledgements do not imply full or partial endorsement by the acknowledged. Further contributions, comments, and criticism are welcome. Please send your remarks to lhc-higgs-properties-convener@cern.ch.

1 Introduction

LHC Higgs Cross Section Working Group is concentrated on various steps of the analysis chain:

Data $\to$ **Pseudo-observables** $\to$ **Model-independent EFT** $\to$ **BSM Models** .

This note concerns model-independent interpretations of the data in the framework of effective field theory (EFT) beyond the Standard Model (SM), which is a part of the scope of the Working Group 2. The purpose of this note is to propose a common EFT language that could be universally used in LHC Higgs analyses.

In the EFT approach, the basic assumption is that the mass scale Λ of new particles beyond the SM is larger than the electroweak scale v, $\Lambda \gg v$. If this is the case, physics at energies $E \ll \Lambda$ can be parametrized by the SM Lagrangian supplemented by a set of higher-dimensional operators. These operators are constructed out of the SM fields, and respect the local $SU(3) \times SU(2) \times U(1)$ symmetry of the SM. The coefficients of $d > 4$-dimensional operators in the EFT Lagrangian are of order $1/\Lambda^{d-4}$, and their contribution to amplitudes of physical processes at the energy scale of order v scales as $(v/\Lambda)^{d-4}$. The leading new physics effects are expected from operators of dimension $d = 6$ whose effects scale as $(v/\Lambda)^2$ (all dimension-5 operators violate the lepton number; experimental constraints dictate that their coefficients must be suppressed at the level unobservable at the LHC). Since $(v/\Lambda)^2 < 1$ by construction, EFT is suitable to describe *small* deviations from the SM predictions, except for observables that vanish or are suppressed by small parameters in the SM.

A *basis* is a complete, non-redundant set of dimension-6 operators. 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 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. In this note we propose a basis that is particularly convenient for LHC Higgs analyses.

Preparing this proposal, we have taken into account the following guidelines.

- Formulation should be simple enough that it can be used by people not acquainted with nuts and bolts of EFTs.
- Relationship between parameters of the EFT and (pseudo)-observables should be transparent.
- Constraints on EFT parameters from electroweak precision observables should be easy to impose.
- The formalism should be simple to implement in monte carlo codes.
- The formalism should be flexible enough, such that, in the future, the application scope may be extended beyond the original one.
- Connection to the pseudo-observables in the *extended kappa formalism* should be straightforward.
- Limits of the EFT validity range should be easy to define.

The basic details of our proposal are the following.

- We restrict to EFT with dimension-6 operators in the *linear* formulation of electroweak symmetry breaking (in other words, the Higgs boson belongs to a doublet of the weak $SU(2)$ group). Non-linear effects may arise in physically motivated cases but, due to a larger number of independent parameters, the corresponding formalism is less predictive and would be more difficult to implement in experimental analyses. Note that this more general case can be constrained by pseudo-observables defined in the extended kappa formalism, which is being developed in parallel by the LHCXSWG2.
- In the spirit of Ref. [1], we proceed with a classification of the operators that more easily map to independent interaction terms of the SM mass eigenstates, in particular of the W, Z, and the Higgs boson. Such interaction terms do not necessarily correspond to $SU(2)$ invariant operators and do not form a complete basis. However, they allow us to identify a set of *independent couplings* from which a complete basis of $SU(2)$-invariant terms is constructed. We denote the latter the *Higgs basis*. The advantage of this formulation is that the effective couplings are related in a simpler way to quantities observable in experiment, compared to other proposals.
- We choose the independent couplings such that the constraints from the Z and W partial decay widths (measured with a per-mille precision by the LEP experiment) can be easily incorporated. These are among the most stringent constraints on EFT parameters, and they have an important impact on possible signals in Higgs searchers. It is unlikely that, at any point in the future, the precision of LHC Higgs searches will be such that the couplings constrained by LEP can be probed

by the LHC with a comparable accuracy. Therefore it is recommended that the the electroweak constraints on Z and W boson couplings to fermions are always imposed when analyzing LHC data, especially on Higgs physics. Other precision observables, such as WW production or off-shell fermion scattering, are numerically less important and we do not discuss the corresponding constraints in this note (see e.g. [2, 3, 4] for a recent discussion).

- The disadvantage of the Higgs basis is that the operator list is cumbersome, being defined by the identification of a set of independent interaction terms after electroweak symmetry breaking. For this reason, we also map the Higgs basis to a set of $SU(3) \times SU(2) \times U(1)$ manifestly invariant operators before electroweak symmetry breaking. For the latter, in this note we use operators in the Warsaw basis of Ref. [5], but it is straightforward to work out a map to any other basis used in the literature. Working with $SU(3) \times SU(2) \times U(1)$ invariant operators may be more convenient for certain calculations (for example, when renormalization group running of the Wilson coefficients needs to be calculated).

- We do not impose flavor universality. While generic constraints on flavor violation are strong, it is plausible that there is a large hierarchy between the coefficients of dimension-6 operators corresponding to different fermion generations. In particular, many models predict the coefficients of operators involving the 3rd generation much larger than those involving the first two generations. Keeping the more general approach will allow us to obtain much more robust constraints on new physics.

- We allow CP violating operators to be present in our basis. In particular, we discuss the most general set of Higgs couplings to matter that include CP violating couplings.

- We assume dimension-6 operators conserve the baryon and lepton number.

In Section 2, to define our notation and conventions, we write down the Standard Model (SM) Lagrangian. In Section 3 we define the Higgs basis, which is the basis we propose for LHC Higgs analyses. The dictionary between the independent couplings and Wilson coefficients of $SU(3) \times SU(2) \times U(1)$ invariant dimension-6 operators in the Warsaw basis is worked out in Section 4.

2 Standard Model Lagrangian

The SM Lagrangian in our notation takes the form

$$\begin{aligned}
\mathcal{L}^{\rm SM} &= -\frac{1}{4}G^a_{\mu\nu}G^a_{\mu\nu} - \frac{1}{4}W^i_{\mu\nu}W^i_{\mu\nu} - \frac{1}{4}B_{\mu\nu}B_{\mu\nu} + D_\mu H^\dagger D_\mu H + \mu_H^2 H^\dagger H - \lambda(H^\dagger H)^2 \\
&+ \sum_{f\in q,\ell} i\bar{f}_L\gamma_\mu D_\mu f_L + \sum_{f\in u,d,e} i\bar{f}_R\gamma_\mu D_\mu f_R \\
&- \left[\tilde{H}^\dagger \bar{u}_R Y^u q_L + H^\dagger \bar{d}_R Y^d V^\dagger_{\rm CKM} q_L + H^\dagger \bar{e}_R Y^\ell \ell_L + \text{h.c.}\right]. \qquad (2.1)
\end{aligned}$$

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 $\tilde{H}_i = \epsilon_{ij} H^*_k$. It acquires the vacuum expectation value (VEV) $\langle H^\dagger H \rangle = v^2/2$ In the unitary gauge, $H = (0, (v+h)/\sqrt{2})$, where h is the Higgs boson field. After electroweak symmetry breaking, the electroweak gauge boson mass eigenstates are defined as $W^\pm = (W^1 \mp i W^2)/\sqrt{2}$, $Z = c_\theta W^3 - s_\theta B$, $A = s_\theta W^3 + c_\theta B$, where $c_\theta = \sqrt{1-s^2_\theta}$. 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 $m_{f_i} = v[Y^f]_{ii}/\sqrt{2}$.

For later convenience, we explicitly write down the gauge boson mass terms:

$$\mathcal{L}^{\rm SM}_{\rm mass} = \frac{g^2 v^2}{4} W^+_\mu W^-_\mu + \frac{(g^2+g'^2)v^2}{8} Z_\mu Z_\mu, \tag{2.2}$$

the gauge boson couplings to fermions:

$$\mathcal{L}^{\rm SM}_{aff} = eA_\mu \sum_{f\in u,d,e} Q_f \bar{f}\gamma_\mu f + g_s G^a_\mu \sum_{f\in u,d} \bar{f}\gamma_\mu T^a f, \tag{2.3}$$

$$\begin{aligned} \mathcal{L}^{\rm SM}_{vff} &= \frac{g}{\sqrt{2}} \left(W^+_\mu \bar{u}_L \gamma_\mu V_{\rm CKM} d_L + W^+_\mu \bar{\nu}_L \gamma_\mu e_L + {\rm h.c.} \right) \\ &+ \sqrt{g^2+g'^2} Z_\mu \sum_{f\in u,d,e,\nu} \left(T^3_f \bar{f}_L \gamma_\mu f_L - s^2_\theta Q_f \bar{f}\gamma_\mu f \right), \end{aligned} \tag{2.4}$$

and the couplings of a single Higgs boson to gauge bosons and fermions:

$$\mathcal{L}^{\rm SM}_h = \frac{h}{v}\left[\frac{g^2 v^2}{2} W^+_\mu W^-_\mu + \frac{(g^2+g'^2)v^2}{4} Z_\mu Z_\mu - \sum_f m_f \bar{f} f \right]. \tag{2.5}$$

These couplings depend on just 4 input parameters: g_s, g, g', and v. These are customarily derived from the observable Fermi constant $G_F = 1/(\sqrt{2}v^2)$ (more precisely, from the measured muon lifetime $\tau_\mu = 384\pi^3 v^4/m^5_\mu$), Z boson mass m_Z, low-energy electromagnetic coupling $\alpha(0) = e^2/4\pi$, and jet production at LEP.

3 Higgs Basis

We present the effective dimension-6 Lagrangian in the linear realization of electroweak symmetry in a formalism inspired by (but not identical to) Ref. [1]. The goal is to choose a particular basis of operators that can be more easily connected (at least at the tree-level) to observable quantities in Higgs physics. The basis, which we call the

Higgs basis, is spanned by a particular combination of dimension-6 operators. Each of these combinations map to a simple interaction term of the SM mass-eigenstate fields that can be probed by experiment. The coefficients multiplying these combinations in the Lagrangian are called the *independent couplings*. In order to make the Higgs basis convenient to study Higgs physics, the couplings of W and Z bosons to fermions and single Higgs couplings to the SM fermions and gauge bosons are chosen among the independent couplings.

We stress that the Higgs basis should be regarded as one of many possible bases of the dimension-6 Lagrangian beyond the SM. In particular, the independent couplings can be related by a linear transformation to parameters defining any other such basis in the literature, for example the Warsaw [5] or to the SILH [6] basis. At the same time, the independent couplings and be easily connected to Higgs *pseudo-observables* at the amplitude level, as defined e.g. in Ref. [7].

In our Lagrangian, by construction, all kinetic terms are canonically normalized, the photon and the gluon interact with fermions as in Eq. (2.3), there is no kinetic mixing between the Z boson and the photon, and there is no correction to the Z boson mass term. These features can be always achieved, without any loss of generality, by using equations of motion, integrating by parts, and redefining the fields and couplings. All the couplings are defined at the scale $\mu = m_h$. In the complete effective Lagrangian each independent coupling multiplies an independent combination of $SU(3) \times SU(2) \times U(1)$ invariant operators (such combinations formally define the operator basis). However, we find it more transparent to define the independent couplings via the interaction terms of SM mass eigenstates in the Lagrangian after electroweak symmetry breaking. See the next section for the expressions of the independent couplings in terms of Wilson coefficients of $SU(3) \times SU(2) \times U(1)$ invariant operators.

We choose the following set of independent couplings relevant for precision pole observables at tree-level:

$$\textbf{EWPT independent couplings}: \quad \delta g_L^{Ze},\ \delta g_R^{Ze},\ \delta g_L^{W\ell},\ \delta g_L^{Zu},\ \delta g_R^{Zu},\ \delta g_L^{Zd},\ \delta g_R^{Zd},\ \delta g_R^{Wq},\ \delta m, \tag{3.1}$$

where all δg are 3×3 matrices in the generation space. These couplings are defined via corrections of the W and Z couplings to the SM fermions in Eq. (2.3) and to the W boson mass in Eq. (2.2):

$$\begin{aligned} \mathcal{L}_{\rm ewpt}^{(1)} &= \frac{g}{\sqrt{2}}\left(W_\mu^+ \bar{u}_R \gamma_\mu \delta g_R^{Wq} d_R + W_\mu^+ \bar{\nu}_L \gamma_\mu \delta g_L^{W\ell} e_L + \text{h.c.}\right) \\ &+ \sqrt{g^2+g'^2} Z_\mu \sum_{f \in u,d,e} \left[\bar{f}_L \gamma_\mu \delta g_L^{Zf} f_L + \bar{f}_R \gamma_\mu \delta g_R^{Zf} f_R\right] \\ &+ 2\delta m \frac{g^2 v^2}{4} W_\mu^+ W_\mu^- \end{aligned} \tag{3.2}$$

These parameters form a complete set to describe all single on-shell Z and W decay and production processes within an EFT with linear realization of electroweak symmetry. The parameters in Eq. (3.1) are free parameters from the effective field theory viewpoint but, as we argue in more detail near the end of this section, they are typically strongly constrained by precision measurements of Z and W production and decays at LEP.

In order to describe the single Higgs production and decay in various channels we need in addition the following set of independent couplings:

$$\textbf{Higgs independent couplings}: \quad \delta c_w,\ \delta c_z,\ c_{gg},\ c_{\gamma\gamma},\ c_{z\gamma},\ c_{zz},\ \tilde{c}_{gg},\ \tilde{c}_{\gamma\gamma},\ \tilde{c}_{z\gamma},\ \tilde{c}_{zz},\ \delta y^u,\ \delta y^d,\ \delta y^e,\ \sin\phi^u,\ \sin\phi^d,\ \sin\phi^\ell. \tag{3.3}$$

These parameters do not affect the pole observables, therefore they are only weakly constrained; typically the strongest limits on these parameters come from Higgs studies at the LHC. The couplings listed in the first line of Eq. (3.3) are defined via the Higgs boson couplings to the SM gauge bosons:

$$\begin{aligned}\mathcal{L}_{\rm hvv}^{(1)} = \frac{h}{v}\Big[& 2\delta c_w m_W^2 W_\mu^+ W_\mu^- + \delta c_z m_Z^2 Z_\mu Z_\mu \\ & + c_{gg}\frac{g_s^2}{4} G_{\mu\nu}^a G_{\mu\nu}^a + c_{\gamma\gamma}\frac{e^2}{4} A_{\mu\nu}A_{\mu\nu} + c_{z\gamma}\frac{eg}{2c_\theta} Z_{\mu\nu}A_{\mu\nu} + c_{zz}\frac{g^2}{4c_\theta^2} Z_{\mu\nu}Z_{\mu\nu} \\ & + \tilde{c}_{gg}\frac{g_s^2}{4} G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a + \tilde{c}_{\gamma\gamma}\frac{e^2}{4} A_{\mu\nu}\tilde{A}_{\mu\nu} + \tilde{c}_{z\gamma}\frac{eg}{2c_\theta} Z_{\mu\nu}\tilde{A}_{\mu\nu} + \tilde{c}_{zz}\frac{g^2}{4c_\theta^2} Z_{\mu\nu}\tilde{Z}_{\mu\nu}\Big].\end{aligned} \tag{3.4}$$

Here $X_{\mu\nu} = \partial_\mu X_\nu - \partial_\nu X_\mu$, and $\tilde{X}_{\mu\nu} = \epsilon_{\mu\nu\rho\sigma}\partial_\rho X_\sigma$. The parameters δy and $\sin\phi$ listed in the second line of Eq. (3.3) are 3×3 real matrices each, defined via the Higgs boson couplings to the SM fermions:

$$\mathcal{L}_{\rm hff} = -\frac{h}{v}\sum_{f\in u,d,e}\sum_{ij}\sqrt{m_{f_i}m_{f_j}}[\delta y^f]_{ij}\left[\cos\phi_{ij}^f \bar{f}_i f_j - i\sin\phi_{ij}^f \bar{f}_i\gamma_5 f_j\right]. \tag{3.5}$$

One could also introduce another set of independent Higgs couplings corresponding to the dipole interaction terms in the Lagrangian of the form $h\bar{f}\sigma_{\mu\nu}fF_{\mu\nu}$, which could affect $h\to 4f$ decays. However, precision measurements (in particular, anomalous magnetic and electric moments of fermions) imply that these couplings must be suppressed at the level that makes them unobservable at the LHC. Moreover, the contribution from the dipole terms does not interfere with the SM amplitudes, which means the corresponding couplings enter at the quadratic level and are therefore suppressed for the size of the couplings within the validity regime of the EFT. Therefore, for simplicity, we do not explicitly write these couplings in this note.

At the level of the dimension-6 Lagrangian, several other Higgs couplings can be expressed by the independent couplings in Eq. (3.1) and Eq. (3.3); we call them the *dependent* couplings. Of course, the choice which couplings are independent and which are dependent is subjective and dictated by convenience; in our case, it is motivated by the fact that the couplings in Eq. (3.1) and Eq. (3.3) are more easily mapped to observables constrained by electroweak precision tests and Higgs searches. One group of dependent couplings are the couplings of the Z boson to neutrinos, and of the W boson to quarks:

$$\mathcal{L}_{\rm ewpt}^{(2)} = \sqrt{g^2+g'^2} Z_\mu \bar{\nu}_L\gamma_\mu \delta g_L^{Z\nu}\nu_L + \left[g W_\mu^+ \bar{u}_i\gamma_\mu \delta g_L^{Wq} V_{\rm CKM} d_L + \text{h.c.}\right], \tag{3.6}$$

which can be expressed by the independent couplings as:

$$\delta g_L^{Z\nu} = \delta g_L^{Ze} + \delta g_L^{W\ell}, \qquad \delta g_L^{Wq} = \delta g_L^{Zu} - \delta g_L^{Zd}. \tag{3.7}$$

Note that we choose the W couplings to leptons (rather than the Z couplings to neutrinos) as our independent couplings, because in the flavor non-universal case the former are more directly constrained by experiment (in particular, in leptonic W decays measured at LEP). Another group of dependent couplings are the coefficients of 2-derivative Higgs boson couplings to W bosons:

$$\mathcal{L}_{\rm hvv}^{(2)} = \frac{h}{v}\left[c_{ww}\frac{g^2}{2}W_{\mu\nu}^+W_{\mu\nu}^- + \tilde{c}_{ww}\frac{g^2}{2}W_{\mu\nu}^+\tilde{W}_{\mu\nu}^-\right], \tag{3.8}$$

where c_{ww} and $\tilde{c}_{ww}$ can be expressed by the independent couplings as

$$\begin{aligned} c_{ww} &= c_{zz} + 2s_\theta^2 c_{z\gamma} + s_\theta^4 c_{\gamma\gamma}, \\ \tilde{c}_{ww} &= \tilde{c}_{zz} + 2s_\theta^2 \tilde{c}_{z\gamma} + s_\theta^4 \tilde{c}_{\gamma\gamma}. \end{aligned} \tag{3.9}$$

Yet another group of dependent couplings are the contact interactions between the Higgs, electroweak gauge bosons, and fermions:

$$\begin{aligned} \mathcal{L}_{hvff} &= \sqrt{2}g\frac{h}{v}W_\mu^+\left(\bar{u}_L\gamma_\mu c_L^{Wq}V_{\rm CKM}d_L + \bar{u}_R\gamma_\mu c_R^{Wq}d_R + \bar{\nu}_L\gamma_\mu c_L^{W\ell}e_L\right) + \text{h.c.} \\ &+ 2\frac{h}{v}\sqrt{g^2+g'^2}Z_\mu\left[\sum_{f=u,d,e,\nu}\bar{f}_L\gamma_\mu c_L^{Zf}f_L + \sum_{f=u,d,e}\bar{f}_R\gamma_\mu c_R^{Zf}f_R\right], \end{aligned} \tag{3.10}$$

where c^{Wf} are related to the independent couplings by

$$\begin{aligned} c^{Zf} &= \delta g^{Zf} + \frac{I_3}{2}\left\{[\delta c_w - \delta c_z - 4\delta m]\left[T_f^3\frac{g^2}{g'^2} + Y_f\right] + T_f^3 g^2\left[c_{\gamma\gamma}\frac{e^2}{g^2+g'^2} + c_{z\gamma}\frac{g^2-g'^2}{g^2+g'^2} - c_{zz}\right]\right\}, \\ c^{Wf} &= \delta g^{Wf} + \frac{I_3}{2}\left\{\frac{g^2}{g'^2}[\delta c_w - \delta c_z - 4\delta m] + g^2\left[c_{\gamma\gamma}\frac{e^2}{g^2+g'^2} + c_{z\gamma}\frac{g^2-g'^2}{g^2+g'^2} - c_{zz}\right]\right\}, \end{aligned} \tag{3.11}$$

and I_3 is the 3×3 unit matrix. Note that, using equations of motion, 3 combinations of the couplings in Eq. (3.10) can be traded for Higgs 2-derivative interactions with gauge bosons that do not involve fermions: $hZ_\mu\partial_\nu Z_{\nu\mu}$, $hZ_\mu\partial_\nu A_{\nu\mu}$, and $hW_\mu^\pm\partial_\nu W_{\nu\mu}^\mp$. For $\delta g^{Vf} = 0$ one can set *all* contact interaction terms to zero, in exchange for these three hVV interaction terms.

The couplings in Eq. (3.1) and Eq. (3.3) are not yet a full set of independent couplings defining the Higgs basis. Similarly, the Lagrangian terms $\mathcal{L}_{\rm ewpt}^{(1,2)}$, $\mathcal{L}_{\rm hff}^{(1)}$, $\mathcal{L}_{\rm hvv}^{(1,2)}$, $\mathcal{L}_{hvff}$ do not correspond to a complete basis of dimension-6 operators. To complete the basis, one needs to add $\mathcal{L}_{\rm other}$ to the Lagrangian with the independent couplings:

$$\textbf{Other independent couplings}: \qquad c_{4f},\ c_{Vf},\ c_{3V},\ c_{3h}, \tag{3.12}$$

where the labels schematically refer to 4-fermion operators, dipole-type operators, gauge boson self-interactions involving 3 field-strength tensors, and the cubic Higgs boson self-interaction. Moreover, $\mathcal{L}_{\rm other}$ contains additional dependent couplings, in particular the ones describing interactions of the SM particles with 2 Higgs bosons. The full Lagrangian in the Higgs basis is given by

$$\mathcal{L}_{\rm Higgs\,Basis} = \mathcal{L}^{\rm SM} + \mathcal{L}_{\rm ewpt}^{(1)} + \mathcal{L}_{\rm ewpt}^{(2)} + \mathcal{L}_{\rm hff} + \mathcal{L}_{\rm hvv}^{(1)} + \mathcal{L}_{\rm hvv}^{(2)} + \mathcal{L}_{hvff} + \mathcal{L}_{\rm other}, \tag{3.13}$$

At the tree level, the terms in $\mathcal{L}_{\rm other}$ are not relevant for electroweak precision tests and single Higgs production and decay. The full expression of $\mathcal{L}_{\rm other}$ will be presented in the future.

In total, the Higgs basis, much as any complete basis at the dimension-6 level, is parametrized by 2499 independent couplings [8]. One should not, however, be intimidated by this number. The point is that a much smaller subset in Eq. (3.3) is adequate for EFT analyses of Higgs data at the leading order in new physics parameters .

A few final comments:

- The relations between independent and dependent couplings in Eq. (3.7), Eq. (3.9), and Eq. (3.11) are consequences of the *linear* realization of electroweak symmetry breaking at the level of dimension-6 EFT operators. *They are an essential part of the definition of the Higgs basis.* If the independent and dependent couplings were unrelated, then $\mathcal{L}_{\rm Higgs\,Basis}$ would not be a dimension-6 basis but would belong to a more general class of theories. Such theories are outside of the scope of this note, however they will be discussed in the framework of the extended kappa formalism.

- The independent couplings in Eq. (3.1) are probed by precision measurements of Z and W production and decays at LEP. In particular, assuming vertex corrections are flavor blind, all the independent couplings in Eq. (3.1) are constrained to be smaller than to $\mathcal{O}(10^{-3})$ (for the leptonic vertex corrections and $\delta m \equiv \delta m_W/m_W$), or $\mathcal{O}(10^{-2})$ (for the quark vertex corrections) [2, 9, 4]. Dropping the assumption of flavor blindness, one cannot constrain all vertex corrections in a completely general way. However, all the leptonic vertex corrections, and bottom and charm quark vertex corrections are still constrained at the level of $\mathcal{O}(10^{-2})$ in a model-independent way by the measurements of leptonic W and Z width and heavy flavor asymmetries. These constraints imply these couplings are too small to have any measurable effects at the LHC, therefore we recommend to impose the electroweak bounds on such constraints before analysing LHC data. The light quark vertex corrections are not constrained in a model independent way; only one combination of them is constrained by measurements of the hadronic Z decays at LEP. Furthermore, the top quark vertex corrections are poorly constrained (at the $\mathcal{O}(1)$ level), by experiment, especially the right-handed top couplings to Z. If feasible, the light quark and top couplings should be considered as free parameters in experimental analyses at the LHC, as this may provide new valuable information to constrain these couplings.

- The Higgs basis is well suited for extracting constraints on dimension-6 operators from experimental data. However, it may not be the optimal basis for some applications. In particular, computing renormalization group running of the couplings or matching to concrete BSM model may be more straightforward in the language of $SU(3) \times SU(2) \times U(1)$ invariant operators.

- Most of the independent Higgs couplings in Eq. (3.3) are probed *only* by Higgs searches. However, 3 linear combinations of these couplings can also be probed by pair production of electroweak gauge bosons (a.k.a, triple gauge couplings measurements). More precisely, in the Higgs basis the deviations of the cubic WWZ

and $WW\gamma$ interactions from the SM in the parametrization of Ref. [10] can be expressed by the independent couplings as

$$\begin{aligned}
\delta g_{1,z} &= -\frac{g^2+g'^2}{2}\left(c_{\gamma\gamma}\frac{e^2}{g^2+g'^2}+c_{z\gamma}\frac{g^2-g'^2}{g^2+g'^2}-c_{zz}\right)+\frac{g^2+g'^2}{2g'^2}\left(\delta c_z-\delta c_w+4\delta m\right),\\
\delta\kappa_\gamma &= -\frac{g^2}{2}\left(c_{\gamma\gamma}\frac{e^2}{g^2+g'^2}+c_{z\gamma}\frac{g^2-g'^2}{g^2+g'^2}-c_{zz}\right),\\
\tilde\kappa_\gamma &= -\frac{g^2}{2}\left(\tilde c_{\gamma\gamma}\frac{e^2}{g^2+g'^2}+\tilde c_{z\gamma}\frac{g^2-g'^2}{g^2+g'^2}-\tilde c_{zz}\right),
\end{aligned}\tag{3.14}$$

and $\delta\kappa_z = \delta g_{1,Z} - t_\theta^2\delta\kappa_\gamma$, $\tilde\kappa_z = -t_\theta^2\tilde\kappa_\gamma$. Note that in the Higgs basis we propose $\delta g_{1,z}$, $\delta\kappa_\gamma$, and $\tilde\kappa_\gamma$ are *dependent* couplings, unlike in Ref. [1].

- Customarily, the SM electroweak parameters are extracted from $\alpha(0)$, m_Z and G_F. One could also use m_W instead of G_F, as suggested in Ref. [2]. This formalism leads to the same relations between the independent and dependent couplings as written down here, except that $\delta m = 0$ by definition. The downside of this formalism is that the SM predictions for all observables would have to be recalculated, as all existing calculations use G_F as an input.

- Additional independent couplings, beyond those in Eq. (3.1) and in Eq. (3.3), are needed to describe, in full generality, the electroweak gauge boson pair production, and the double Higgs production. This is left for (near) future work of the LHCXSWG2.

- The number of independent couplings in Eq. (3.3) relevant for Higgs observables is still large. At the early stages of the LHC run-2 it may be reasonable to employ simplified analyses with a smaller number of parameters. There are several motivated assumptions about underlying new physics theory that reduce the number of parameters:

  - *Flavor blindness*, in which case the matrices δy^f and $\sin\phi^f$ reduce to a single number for each $f = u, d, e$.
  - *CP conservation*, in which case $\tilde c_i = \sin\phi^f = 0$.
  - *Custodial symmetry* or *universalness.* The constraints in this case will be written down later.

  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 should be consistently imposed, so as to preserve the $SU(2)_L$ local symmetry.

- The independent couplings are formally of order v^2/Λ^2, where Λ is the scale of new physics. For completeness, it is important to define the range of independent couplings such that the EFT description is valid. The rule of thumb is that this is the case for $|c_i| \lesssim 1$; more sophisticated criteria will be worked out in the future when specific Higgs processes are discussed.

4 Map to Dimension-6

We turn to discussing the map between the independent couplings introduced in Section 3 and coefficients of dimension-6 operators in the electroweak basis before electroweak symmetry breaking. The complete set of dimension-6 operators can be written in many different equivalent bases which are related by the use of equations of motion and integration by parts. Here we work with the so-called *Warsaw basis* of Ref. [5, 8], which is distinguished by the simplest tensor structure of the higher-dimensional operators.

The Lagrangian in the Warsaw basis is given by[1]

$$\mathcal{L}_{\rm warsaw} = \mathcal{L}^{\rm SM} + \frac{1}{v^2}\sum_i c_i O_i, \tag{4.1}$$

where the SM Lagrangian $\mathcal{L}^{\rm SM}$ was introduced in Section 2, and the dimension-6 operators O_i are summarized in Table 1.

To map the coefficients of dimension-6 operators into the independent couplings in Eq. (3.1) and Eq. (3.3) , we need first to bring $\mathcal{L}_{\rm warsaw}$ into the same form as $\mathcal{L}_{\rm Higgs\ Basis}$ in Eq. (3.13). This can be achieved by a series of transformations using equations of motion, integration by parts, and rescaling of the fields and couplings. To begin with, the operator O_{WB} leads to a kinetic mixing between the hypercharge and SU(2) gauge bosons, $O_{WB} \to -1/2 g g' W^3_{\mu\nu} B_{\mu\nu}$. To get rid of it, we use the equations of motion:

$$\begin{aligned}
\partial_\nu B_{\nu\mu} &= g' \frac{(v+h)^2}{4}\left(g W^3_\mu - g' B_\mu\right) - g' j^Y_\mu, \\
\partial_\nu W^3_{\nu\mu} &= -g \frac{(v+h)^2}{4}\left(g W^3_\mu - g' B_\mu\right) - g j^3_\mu - g\epsilon^{3jk} W^j_\nu W^k_{\nu\mu},
\end{aligned} \tag{4.2}$$

where $j^Y_\mu = \sum_f Y_f \bar f \gamma_\mu f$, and $j^3_\mu = \bar q \gamma_\mu T^3 P_L q + \bar\ell \gamma_\mu T^3 P_L \ell$. Using this,

$$\begin{aligned}
c_{WB}\frac{g g'}{2} W^3_{\mu\nu} B_{\mu\nu} &\to c_{WB} e^2 \left[\frac{(v+h)^2}{4}\left(g W^3_\mu - g' B_\mu\right)^2 - g W^3_\mu j^Y_\mu - g' B_\mu j^3_\mu\right] + \ldots \\
&= c_{WB} e^2 \left[\frac{(g^2+g'^2)(v+h)^2}{4} Z^2_\mu - e A_\mu j^{\rm em}_\mu + \sqrt{g^2+g'^2} Z_\mu \left(j^3_\mu - c^2_\theta j^{\rm em}_\mu\right)\right] + \ldots,
\end{aligned} \tag{4.3}$$

where $j^{\rm em}_\mu = j^3_\mu + j^Y_\mu$ is the electromagnetic current, and dots stand for cubic gauge boson intereactions that are not relevant for the present discussion

Next, the operators O_H, O_{BB}, O_{WW}, and O_{GG} change the normalization of the kinetic terms of the Higgs and gauge bosons. To recover the canonical normalization, we redefine the fields as

$$h \to h\,(1 - c_H)\,, \; B_\mu \to B_\mu \left(1 + \frac{c_{BB} g'^2}{4}\right), \; W^i_\mu \to W^i_\mu \left(1 + \frac{c_{WW} g^2}{4}\right), \; G^a_\mu \to G^a_\mu \left(1 + \frac{c_{GG} g_s^2}{4}\right). \tag{4.4}$$

[1] We use a different notation than the original reference. We also replaced the operator $|H^\dagger D_\mu H|^2$ by $(H^\dagger D_\mu H - D_\mu H^\dagger H)^2$. For Yukawa-type operators O_f we subtracted v^2 so that these operators do not contribute to off-diagonal mass terms. This way we avoid tedious rotations of the fermion fields to bring them back to the mass eigenstate basis. 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. (2.1) and Table 1 by defining $f'_{L,R} = U_{L,R} f_{L,R}$, $c_f = U^\dagger_R c'_f U_L$, $Y_f = U^\dagger_R (Y'_f + c'_f/2) U_L$, where $U_{L,R}$ are unitary rotations to the mass eigenstate basis.

We ignore here the contribution of the operator $\tilde{O}_{GG}$ to the OCD θ-term (we can always assume it cancels agains the θ-term in the SM Lagrangian, or is dynamically removed by an axion field).

Finally, we have to make sure that the gauge couplings and the Higgs VEV have the same meaning as in the SM. This is a non-trivial requirement, because dimension-6 operators affect the observables used to extract these parameters. We have seen that the operator O_{WB} shifts the electric charge and the Z boson mass. Similarly, the operator O_T shifts the Z boson mass term. Furthermore, a combination of the $O_{\ell\ell}$ operators leads, via Fierz transformations, to the the 4-fermion coupling $-4c'_{\ell\ell}(\bar{\nu}_{\mu,L}\gamma_\rho\mu_L)(\bar{e}_L\gamma_\rho\nu_{e,L})$ that shifts the Fermi constant. Finally, the leptonic vertex operator $O_{H\ell}$ also shifts the Fermi constant. To undo these effects, we need to ensure that the photon and the gluon couple to the electromagnetic and strong currents as in Eq. (2.3). Furthermore, the Z boson mass term in the Lagrangian should be as in Eq. (2.2), and the tree-level $\mu \to e\bar{\nu}_e\nu_\mu$ decay width should be given by $\Gamma = \frac{m_\mu^5}{384\pi^3 v^4}$. This is achieved by the following redefinition of the coupling constants and the VEV:

$$
\begin{aligned}
g_s &\to g_s\left(1 - c_{GG}\frac{g_s^2}{4}\right), \\
g &\to g\left(1 - c_{WW}\frac{g^2}{4} - c_{WB}\frac{g^2 g'^2}{g^2 - g'^2} + (c_T - \delta v)\frac{g^2}{g^2 - g'^2}\right), \\
g' &\to g'\left(1 - c_{BB}\frac{g'^2}{4} + c_{WB}\frac{g^2 g'^2}{g^2 - g'^2} - (c_T - \delta v)\frac{g'^2}{g^2 - g'^2}\right), \\
v &\to v\,(1 + \delta v),
\end{aligned} \tag{4.5}
$$

where $\delta v = ([c'_{H\ell}]_{11} + [c'_{H\ell}]_{22})/2 - c'_{\ell\ell}$.

After the transformations in Eq. (4.3), Eq. (4.4), Eq. (4.5), the Lagrangian takes the same form as $\mathcal{L}_{\rm Higgs\,Basis}$. The dictionary between the coefficients of dimension-6 operators and the couplings in $\mathcal{L}_{\rm Higgs\,Basis}$ goes as follows. The shift of the W boson mass is given by

$$
\delta m = \frac{1}{g^2 - g'^2}\left[-g^2 g'^2 c_{WB} + g^2 c_T - g'^2 \delta v\right] \tag{4.6}
$$

The shift of W and Z boson couplings to leptons are given by

$$
\begin{aligned}
\delta g_L^{W\ell} &= c'_{H\ell} + f(1/2, 0) - f(-1/2, -1), \\
\delta g_L^{Z\nu} &= \frac{1}{2}c'_{H\ell} - \frac{1}{2}c_{H\ell} + f(1/2, 0), \\
\delta g_L^{Ze} &= -\frac{1}{2}c'_{H\ell} - \frac{1}{2}c_{H\ell} + f(-1/2, -1), \\
\delta g_R^{Ze} &= -\frac{1}{2}c_{He} + f(0, -1),
\end{aligned} \tag{4.7}
$$

where

$$
f(T^3, Q) = I_3\left[-Q c_{WB}\frac{g^2 g'^2}{g^2 - g'^2} + (c_T - \delta v)\left(T^3 + Q\frac{g'^2}{g^2 - g'^2}\right)\right], \tag{4.8}
$$

and I_3 is the 3×3 identity matrix. The shifts of W and Z boson couplings to quarks are given by

$$\begin{aligned}
\delta g_L^{Wq} &= c'_{Hq} + f(1/2, 2/3) - f(-1/2, -1/3), \\
\delta g_R^{Wq} &= -\frac{1}{2} c_{Hud}, \\
\delta g_L^{Zu} &= \frac{1}{2} c'_{Hq} - \frac{1}{2} c_{Hq} + f(1/2, 2/3), \\
\delta g_L^{Zd} &= -\frac{1}{2} c'_{Hq} - \frac{1}{2} c_{Hq} + f(-1/2, -1/3), \\
\delta g_R^{Zu} &= -\frac{1}{2} c_{Hu} + f(0, 2/3), \\
\delta g_R^{Zd} &= -\frac{1}{2} c_{Hd} + f(0, -1/3).
\end{aligned} \tag{4.9}$$

The shifts of the Higgs couplings to W and Z are given by

$$\begin{aligned}
\delta c_w &= -c_H - c_{WB} \frac{2g^2 g'^2}{g^2 - g'^2} + 2c_T \frac{g^2}{g^2 - g'^2} - \delta v \frac{g^2 + g'^2}{g^2 - g'^2}, \\
\delta c_z &= -c_H - 2c_T - \delta v.
\end{aligned} \tag{4.10}$$

The two-derivative Higgs couplings to gauge bosons are given by

$$\begin{aligned}
c_{gg} &= c_{GG}, \\
c_{\gamma\gamma} &= c_{WW} + c_{BB} - 4c_{WB}, \\
c_{zz} &= \frac{g^4 c_{WW} + g'^4 c_{BB} + 4g^2 g'^2 c_{WB}}{(g^2 + g'^2)^2}, \\
c_{z\gamma} &= \frac{g^2 c_{WW} - g'^2 c_{BB} - 2(g^2 - g'^2) c_{WB}}{g^2 + g'^2}, \\
c_{ww} &= c_{WW}.
\end{aligned} \tag{4.11}$$

The Yukawa interactions are given by

$$\begin{aligned}
\sqrt{m_i m_j} \delta y_{ij}^f \cos\phi_{ij}^f &= \frac{v}{2} \left[c_f + c_f^\dagger \right]_{ij} - m_{f_i} \delta_{ij} \left(c_H + \delta v \right) \\
\sqrt{m_i m_j} \delta y_{ij}^f \sin\phi_{ij}^f &= \frac{v}{2i} \left[c_f - c_f^\dagger \right]_{ij}.
\end{aligned} \tag{4.12}$$

Finally, the contact interactions between a Higgs, a gauge boson, and 2 fermions are given by

$$c_L^{W\ell} = c'_{H\ell}, \quad c_L^{Wq} = c'_{Hq}, \quad c_R^{Wq} = -\frac{1}{2} c'_{Hud}, \tag{4.13}$$

$$c_L^{Z\nu} = \frac{1}{2} c'_{H\ell} - \frac{1}{2} c_{H\ell}, \quad c_L^{Ze} = -\frac{1}{2} c'_{H\ell} - \frac{1}{2} c_{H\ell}, \quad c_R^{Ze} = -\frac{1}{2} c_{He}, \tag{4.14}$$

$$c_L^{Zu} = \frac{1}{2} c'_{Hq} - \frac{1}{2} c_{Hq}, \quad c_L^{Zd} = -\frac{1}{2} c'_{Hq} - \frac{1}{2} c_{Hq}, \quad c_R^{Zu} = -\frac{1}{2} c_{Hu}, \quad c_R^{Zd} = -\frac{1}{2} c_{Hd}. \tag{4.15}$$

From these expression one can derive the relations between dependent and independent couplings listed in the previous section.

To summarize, in the Warsaw basis the parameters affecting electroweak precision tests and single Higgs production and decay are the following

$$\begin{gathered} c_H,\ c_T,\ c_{GG},\ c_{WW},\ c_{BB},\ c_{WB},\ \tilde{c}_{GG},\ \tilde{c}_{WW},\ \tilde{c}_{BB},\ \tilde{c}_{WB},\ c'_{\ell\ell}, \\ c'_{H\ell},\ c_{H\ell},\ c_{He},\ c'_{Hq},\ c_{Hq},\ c_{Hu},\ c_{Hd},\ c_{Hud} \\ c^u,\ c^d,\ c^\ell. \end{gathered} \tag{4.16}$$

The linear transformation between these parameters and the independent couplings in Eq. (3.1) and Eq. (3.3) is given in Eqs. (4.6)-(4.15). In principle, one can also perform the LHC analyses in the Warsaw (or any other) basis. One difficulty is that the electroweak precision constraints, which are transparent in the Higgs basis, constrain rather complicated combinations of the parameters in Eq. (4.16). Alternatively, the constraints derived in the Higgs basis can be easily recast into constraints in the Warsaw basis using the map Eqs. (4.6)-(4.15), provided that the former are given with the full correlation matrix.

A similar procedure can be applied to any other basis, for example the SILH one [6]. The relations between the Higgs couplings and the Wilson coefficients in the SILH basis can be found e.g. in Refs. [11, 12]. The explicit transformation between the Higgs and SILH bases will be added in the future.

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H^4D^2 and H^6

O_H	$\left[\partial_\mu(H^\dagger H)\right]^2$
O_T	$\left(H^\dagger \overleftrightarrow{D}_\mu H\right)^2$
O_{6H}	$(H^\dagger H)^3$

f^2H^3

O_e	$-(H^\dagger H - \frac{v^2}{2})\bar{e}H^\dagger \ell$
O_u	$-(H^\dagger H - \frac{v^2}{2})\bar{u}\widetilde{H}^\dagger q$
O_d	$-(H^\dagger H - \frac{v^2}{2})\bar{d}H^\dagger q$

V^3D^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}	$g^3 \epsilon^{ijk} W^i_{\mu\nu} W^j_{\nu\rho} W^k_{\rho\mu}$
$O_{3\widetilde{W}}$	$g^3 \epsilon^{ijk} \widetilde{W}^i_{\mu\nu} W^j_{\nu\rho} W^k_{\rho\mu}$

V^2H^2

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}	$g g' H^\dagger \sigma^i H\, W^i_{\mu\nu} B_{\mu\nu}$
$O_{\widetilde{WB}}$	$g g' H^\dagger \sigma^i H\, \widetilde{W}^i_{\mu\nu} B_{\mu\nu}$

f^2H^2D

$O_{H\ell}$	$i\bar{\ell}\gamma_\mu \ell H^\dagger \overleftrightarrow{D}_\mu H$
$O'_{H\ell}$	$i\bar{\ell}\sigma^i\gamma_\mu \ell H^\dagger \sigma^i \overleftrightarrow{D}_\mu H$
O_{He}	$i\bar{e}\gamma_\mu \bar{e} H^\dagger \overleftrightarrow{D}_\mu H$
O_{Hq}	$i\bar{q}\gamma_\mu q H^\dagger \overleftrightarrow{D}_\mu H$
O'_{Hq}	$i\bar{q}\sigma^i\gamma_\mu q H^\dagger \sigma^i \overleftrightarrow{D}_\mu H$
O_{Hu}	$i\bar{u}\gamma_\mu u H^\dagger \overleftrightarrow{D}_\mu H$
O_{Hd}	$i\bar{d}\gamma_\mu d H^\dagger \overleftrightarrow{D}_\mu H$
O_{Hud}	$i\bar{u}\gamma_\mu d \tilde{H}^\dagger D_\mu H$

f^2VHD

O_{eW}	$g\bar{\ell}\sigma_{\mu\nu} e \sigma^i H W^i_{\mu\nu}$
O_{eB}	$g'\bar{\ell}\sigma_{\mu\nu} e H B_{\mu\nu}$
O_{uG}	$g_s\bar{q}\sigma_{\mu\nu} T^a u \widetilde{H}\, G^a_{\mu\nu}$
O_{uW}	$g\bar{q}\sigma_{\mu\nu} u \sigma^i \widetilde{H}\, W^i_{\mu\nu}$
O_{uB}	$g'\bar{q}\sigma_{\mu\nu} u \widetilde{H}\, B_{\mu\nu}$
O_{dG}	$g_s\bar{q}\sigma_{\mu\nu} T^a d H\, G^a_{\mu\nu}$
O_{dW}	$g\bar{q}\sigma_{\mu\nu} d \sigma^i H\, W^i_{\mu\nu}$
O_{dB}	$g'\bar{q}\sigma_{\mu\nu} d H\, B_{\mu\nu}$

$(\bar{L}L)(\bar{L}L)$ and $(\bar{L}R)(\bar{L}R)$

$O_{\ell\ell}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{\ell}\gamma_\mu \ell)$
O_{qq}	$(\bar{q}\gamma_\mu q)(\bar{q}\gamma_\mu q)$
O'_{qq}	$(\bar{q}\gamma_\mu \sigma^i q)(\bar{q}\gamma_\mu \sigma^i q)$
$O_{\ell q}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{q}\gamma_\mu q)$
$O'_{\ell q}$	$(\bar{\ell}\gamma_\mu \sigma^i \ell)(\bar{q}\gamma_\mu \sigma^i q)$
O_{quqd}	$(\bar{q}^j u)\epsilon_{jk}(\bar{q}^k d)$
O'_{quqd}	$(\bar{q}^j T^a u)\epsilon_{jk}(\bar{q}^k T^a d)$
$O_{\ell equ}$	$(\bar{\ell}^j e)\epsilon_{jk}(\bar{q}^k u)$
$O'_{\ell equ}$	$(\bar{\ell}^j \sigma_{\mu\nu} e)\epsilon_{jk}(\bar{q}^k \sigma^{\mu\nu} u)$
$O_{\ell edq}$	$(\bar{\ell}^j e)(\bar{d} q^j)$

$(\bar{R}R)(\bar{R}R)$

O_{ee}	$(\bar{e}\gamma_\mu e)(\bar{e}\gamma_\mu e)$
O_{uu}	$(\bar{u}\gamma_\mu u)(\bar{u}\gamma_\mu u)$
O_{dd}	$(\bar{d}\gamma_\mu d)(\bar{d}\gamma_\mu d)$
O_{eu}	$(\bar{e}\gamma_\mu e)(\bar{u}\gamma_\mu u)$
O_{ed}	$(\bar{e}\gamma_\mu e)(\bar{d}\gamma_\mu d)$
O_{ud}	$(\bar{u}\gamma_\mu u)(\bar{d}\gamma_\mu d)$
O'_{ud}	$(\bar{u}\gamma_\mu T^a u)(\bar{d}\gamma_\mu T^a d)$

$(\bar{L}L)(\bar{R}R)$

$O_{\ell e}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{e}\gamma_\mu e)$
$O_{\ell u}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{u}\gamma_\mu u)$
$O_{\ell d}$	$(\bar{\ell}\gamma_\mu \ell)(\bar{d}\gamma_\mu d)$
O_{qe}	$(\bar{q}\gamma_\mu q)(\bar{e}\gamma_\mu e)$
O_{qu}	$(\bar{q}\gamma_\mu q)(\bar{u}\gamma_\mu u)$
O'_{qu}	$(\bar{q}\gamma_\mu T^a q)(\bar{u}\gamma_\mu T^a u)$
O_{qd}	$(\bar{q}\gamma_\mu q)(\bar{d}\gamma_\mu d)$
O'_{qd}	$(\bar{q}\gamma_\mu T^a q)(\bar{d}\gamma_\mu T^a d)$

Table 1: A complete, non-redundant set of baryon-and-lepton-number-conserving dimension-6 operators built from SM fields [5]. 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. Including the flavor structure and complex conjugates, this table contains 2499 distinct operators [8].