BSM Primaries: the Physics of the SM Effective Field Theory

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The phenomenological implications of the Standard Model (SM) are governed by the accidental symmetry structure of the dimension-4 Lagrangian. In this talk I discuss the next order in an expansion in fields and in derivatives, that parametrize the largest effects of heavy physics beyond the SM. The remaining symmetries of this dimension-6 Lagrangian imply relations between experimental observables that should be used to test the consistency of deviations from the SM, to design new physics searches and to make them more sensitive.

1 Motivation

The Higgs bosons discovery marks the culmination of searches for the Standard Model (SM) of particle physics. All of the SM sectors have finally been probed and most of its parameters accessed experimentally, with different levels of precision. At the same time direct searches for physics beyond the SM (BSM) have been unsuccessful, suggesting the existence of a mass gap between the SM states and any possible mass scale characteristic of the new physics sector.

In this situation, where the energy of our experiments seems to be insufficient to produce BSM degrees of freedom on-shell, we can still hope that their virtual exchange induces some visible effects, as modifications of the interactions between SM states. This represents the main motivation to perform SM precision tests.

For these tests to bear any quantitative physical significance and for their results to be readily interpretable in the framework of searches for new physics, an appropriate parametrization of the possible departures from the SM is necessary. This parametrization is naturally provided by an SM effective field theory (EFT), which groups all possible interaction among the SM fields in a series expansion in inverse powers of the scale of new physics Λ: $\mathcal{L}_{\rm eff} = \mathcal{L}_4 + \mathcal{L}_6 + \cdots$, where $\mathcal{L}_4$ is made of dimension-4 operators and defines what we call the SM Lagrangian, while $\mathcal{L}_6$, that contains dimension-6 operators suppressed by Λ^2, gives the leading BSM effects.[a] From a bottom-up perspective, these interactions can be considered necessary and their coefficients (the scales associated with each of them) can be fixed only through experiments, in the same way as one fixes the SM input parameters through precise measurements of the input observables $(\alpha, m_Z, G_F, \ldots)$. From a top-bottom perspective, on the other hand, specific BSM models can be matched straightforwardly to the EFT description, by integrating out the relevant massive particles. This twofold

[a] Assuming lepton and baryon number conservation.

interpretation of the EFT parametrization, makes it a suitable tool to characterize departures from the SM in such a way that precision SM tests can be turned into searching tools and their results compared with other direct or indirect searches.

In this note I review the leading departures from the SM in an EFT description. Interestingly, of the many accidental symmetries and relations that define the SM Lagrangian, some resist at the leading order in the EFT expansion (equivalently: the number of observables affected by the leading EFT effects, is smaller than the number of operators characterizing the leading EFT Lagrangian). For this reason the EFT analysis implies some relations between observables (the analog of, e.g., the SM relation $m_W = m_Z/\cos\theta_W$), that represent an important piece of information about the BSM structure. In fact, these relations can be used the test the assumptions behind the EFT (e.g. a separation of scales or the exactness of the SM symmetries); alternatively, they can be used to identify the directions that have been weakly probed by current and past experiments and understand which observables deserve particular attention.

2 BSM Primaries

There are several possibilities to write $\mathcal{L}_6$. From a top-bottom perspective, different operator bases for $\mathcal{L}_6$ can facilitate the comparison with explicit BSM scenarios. For instance, the SILH basis [1] was constructed to capture the effects of universal theories (where the new physics couples only to bosons), such as SUSY or Composite Higgs[b], while the basis of Ref. [4] makes the matching with theories with (partially) composite fermions more straightforward. From a bottom-up perspective, however, these formulations are all equivalent as one is only interested in complete sets of operators. In fact, from this point of view, we can treat $\mathcal{L}_6$ in exact analogy with the SM Lagrangian $\mathcal{L}_4$: we chose the most precise experiments to fix its parameters (for the SM, $\mathcal{L}_4$, we typically take $\alpha, m_Z, G_F,\ldots$) and then express all other observables in terms of these *input* parameters (observables in terms of observables). To this end, we must identify a set of well-measured input observables (which are actually affected by the modifications implied in $\mathcal{L}_6$) that allows us to fix the parameters in $\mathcal{L}_6$. This matching between coefficients in $\mathcal{L}_6$ and well-measured observables was performed in Refs. [5,6,7] and named *BSM Primaries* basis, and I summarize it here.

The first important step is to recognize that there is a class of BSM operators, which in the gauge eigenstate basis corresponds to operators of the form $|H|^2 \times \mathcal{L}_{SM}$, which can only be tested in Higgs physics [6,10]. In fact, when these operators are measured in the vacuum $\langle \hat{h} \rangle = v$, they can be absorbed into a redefinition of some SM parameter and they have, therefore, no physical effect. The number of such operators equals the number of SM parameters which, if we limit ourselves to CP conserving quantities and a diagonal flavor structure[c], reduces to eight, which we write as:

$$\Delta\mathcal{L}^h_{\gamma\gamma} = \kappa_{\gamma\gamma}\left(\frac{h}{v}+\frac{h^2}{2v^2}\right)\left[A_{\mu\nu}A^{\mu\nu}+Z_{\mu\nu}Z^{\mu\nu}+2W^+_{\mu\nu}W^{-\mu\nu}\right], \tag{1}$$

$$\Delta\mathcal{L}^h_{Z\gamma} = \kappa_{Z\gamma}\left(\frac{h}{v}+\frac{h^2}{2v^2}\right)\left[t_{\theta_W}A_{\mu\nu}Z^{\mu\nu}+\frac{c_{2\theta_W}}{2c^2_{\theta_W}}Z_{\mu\nu}Z^{\mu\nu}+W^+_{\mu\nu}W^{-\mu\nu}\right], \tag{2}$$

$$\Delta\mathcal{L}^h_{GG} = \kappa_{GG}\left(\frac{h}{v}+\frac{h^2}{2v^2}\right)G^A_{\mu\nu}G^{A\,\mu\nu}, \tag{3}$$

$$\Delta\mathcal{L}^h_{ff} = \delta g^h_{ff}\left(h\bar{f}_L f_R+\text{h.c.}\right)\left(1+\frac{3h}{2v}+\frac{h^2}{2v^2}\right), \tag{4}$$

$$\Delta\mathcal{L}_{3h} = \delta g_{3h}\, h^3\left(1+\frac{3h}{2v}+\frac{3h^2}{4v^2}+\frac{h^3}{8v^3}\right), \tag{5}$$

$$\Delta\mathcal{L}^h_{VV} = \delta g^h_{VV}\left[h\left(W^{+\mu}W^-_\mu+\frac{Z^\mu Z_\mu}{2c^2_{\theta_W}}\right)+\Delta_h\right], \tag{6}$$

[b] See Refs.[2,3], and references therein, for analysis of the contributions to the EFT of Supersymmetric and Composite Higgs models [?,?]. .

[c] These arguments can be easily extended to higher-order effects in a Minimal Flavor Violation (MFV[11]) expansion[6] or to more complicated flavor structures[12].

where $f = u, d, e$ runs over different types of fermion. Here, I denote $\hat{h} \equiv v + h(x)$ the Higgs field, where h denotes the physical Higgs degree of freedom; Δ_h includes interaction which are irrelevant for experiments in the near future[5] and I define $Z_{\mu\nu} \equiv \hat{Z}_{\mu\nu} - igc_{\theta_W} W^+_{[\mu} W^-_{\nu]}$, $A_{\mu\nu} \equiv \hat{A}_{\mu\nu} - igs_{\theta_W} W^+_{[\mu} W^-_{\nu]}$ and $W^{\pm}_{\mu\nu} \equiv \hat{W}^{\pm}_{\mu\nu} \pm igW^{\pm}_{[\mu}(s_{\theta_W} A + c_{\theta_W} Z)_{\nu]}$. Written in this way, $\mathcal{L}_6$ is automatically ready to incorporate the experimental information from measurements of the Higgs decay and production rates: measurements of the rates $h \to \gamma Z, \gamma\gamma, \bar{f}f$, the production channel $GG \to h$ and the – custodial preserving – $hV_\mu V^\mu$ and h^3 vertices (the latter not yet accessible), allow to fix the parameters $\{\kappa_{\gamma Z}, \kappa_{\gamma\gamma}, \delta g^h_{uu}, \delta g^h_{dd}, \delta g^h_{ee}, \kappa_{GG}, \delta g^h_{VV}, \delta g_{3h}\}$. Notice that in the above expressions, and throughout this note, I absorb powers of m^2_W/Λ^2_i or v^2/Λ^2 into the coefficients κ_i and δg_i: in this way, for the EFT to make sense, we expect $\kappa_i, \delta g_i \ll 1$. Unfortunately (see e.g. Refs. [8,9,13]) this is not the case for the δg_i couplings at present, implying that the use of the EFT parametrization in this case is not yet justified. However, for the κ_i couplings the constraints are already very stringent: at the 95% C.L.[8,13,6] $\kappa_{\gamma\gamma} \in [-1.3, 1.8] \times 10^{-3}$, $\kappa_{Z\gamma} \in [-2, 4] \times 10^{-2}$, $\kappa_{GG} \in [-1, 1] \times 10^{-2}$. Then, Eqs. (1-6) automatically imply a prediction that the coefficient of structures like $hZ_{\mu\nu}Z^{\mu\nu}$, which modifies the differential distribution of $h \to ZZ^*$ (see below), receives contributions from Eq. (1) and Eq. (2), but this contribution is limited within the range of given above. A second class of BSM effects contained in $\mathcal{L}_6$ can instead be measured both in Higgs physics and in the vacuum. In the language of effective operators, these effects are associated with structures like $H^\dagger \sigma^a H$, which transforms non-trivially under $SU(2)_L$ and implies measurable EWSB effects for $\langle \hat{h} \rangle = v$. In this case, at present, the measurement of these effects is more easily performed in the vacuum[d] and, for this reason, we parametrize this sector of the effective Lagrangian as,

$$\begin{aligned}
\Delta\mathcal{L}^V_{ee} &= \boldsymbol{\delta g^Z_{eR}} \frac{\hat{h}^2}{v^2} Z^\mu \bar{e}_R \gamma_\mu e_R + \boldsymbol{\delta g^Z_{eL}} \frac{\hat{h}^2}{v^2} \left[Z^\mu \bar{e}_L \gamma_\mu e_L - \frac{c_{\theta_W}}{\sqrt{2}} (W^{+\mu} \bar{\nu}_L \gamma_\mu e_L + \text{h.c.}) \right] \\
&\quad + \boldsymbol{\delta g^Z_{\nu L}} \frac{\hat{h}^2}{v^2} \left[Z^\mu \bar{\nu}_L \gamma_\mu \nu_L + \frac{c_{\theta_W}}{\sqrt{2}} (W^{+\mu} \bar{\nu}_L \gamma_\mu e_L + \text{h.c.}) \right], \qquad (7) \\
\Delta\mathcal{L}^V_{qq} &= \boldsymbol{\delta g^Z_{uR}} \frac{\hat{h}^2}{v^2} Z^\mu \bar{u}_R \gamma_\mu u_R + \boldsymbol{\delta g^Z_{dR}} \frac{\hat{h}^2}{v^2} Z^\mu \bar{d}_R \gamma_\mu d_R + \boldsymbol{\delta g^Z_{dL}} \frac{\hat{h}^2}{v^2} \left[Z^\mu \bar{d}_L \gamma_\mu d_L - \frac{c_{\theta_W}}{\sqrt{2}} (W^{+\mu} \bar{u}_L \gamma_\mu d_L + \text{h.c.}) \right] \\
&\quad + \boldsymbol{\delta g^Z_{uL}} \frac{\hat{h}^2}{v^2} \left[Z^\mu \bar{u}_L \gamma_\mu u_L + \frac{c_{\theta_W}}{\sqrt{2}} (W^{+\mu} \bar{u}_L \gamma_\mu d_L + \text{h.c.}) \right] \qquad (8) \\
\Delta\mathcal{L}_{g^Z_1} &= \boldsymbol{\delta g^Z_1} \Bigg[igc_{\theta_W} \Big(Z^\mu (W^{+\nu} W^-_{\mu\nu} - \text{h.c.}) + Z^{\mu\nu} W^+_\mu W^-_\nu \Big) \\
&\quad - 2gc^2_{\theta_W} \frac{h}{v} \left(W^-_\mu J^\mu_- + \text{h.c.} + \frac{c_{2\theta_W}}{c^3_{\theta_W}} Z_\mu J^\mu_Z + \frac{2s^2_{\theta_W}}{c_{\theta_W}} Z_\mu J^\mu_{em} \right) \left(1 + \frac{h}{2v} \right) \qquad (9) \\
&\quad + \frac{e^2 v}{2c^2_{\theta_W}} h Z_\mu Z^\mu + g^2 c^2_{\theta_W} v\, \Delta_h - g^2 c^2_{\theta_W} \Big(W^+_\mu W^{-\mu} + \frac{c_{2\theta_W}}{2c^4_{\theta_W}} Z_\mu Z^\mu \Big) \Big(\frac{5h^2}{2} + \frac{2h^3}{v} + \frac{h^4}{2v^2} \Big) \Bigg] \\
\Delta\mathcal{L}_{\kappa_\gamma} &= \frac{\boldsymbol{\delta\kappa_\gamma}}{v^2} \Big[ie\hat{h}^2 (A_{\mu\nu} - t_{\theta_W} Z_{\mu\nu}) W^{+\mu} W^{-\nu} + Z_\nu \partial_\mu \hat{h}^2 (t_{\theta_W} A^{\mu\nu} - t^2_{\theta_W} Z^{\mu\nu}) \\
&\quad + \frac{(\hat{h}^2 - v^2)}{2} \Big(t_{\theta_W} Z_{\mu\nu} A^{\mu\nu} + \frac{c_{2\theta_W}}{2c^2_{\theta_W}} Z_{\mu\nu} Z^{\mu\nu} + W^+_{\mu\nu} W^{-\mu\nu} \Big) \Big], \qquad (10)
\end{aligned}$$

all these effects, in the vacuum, can be measured as modifications of SM couplings (meaning that their contribution interferes with the SM in the amplitude-squared) and from a comparison with LEP1 data, we find[15]

$$\begin{aligned}
\boldsymbol{\delta g^Z_{eL}} = 0.4^{\pm 0.5} \times 10^{-3}, \quad \boldsymbol{\delta g^Z_{eR}} = -0.1^{\pm 0.3} \times 10^{-3}, \quad \boldsymbol{\delta g^Z_{\nu}} = -1.6^{\pm 0.8} \times 10^{-3}, \qquad (11) \\
\boldsymbol{\delta g^Z_{uL}} = -2.6^{\pm 1.6} \times 10^{-3}, \quad \boldsymbol{\delta g^Z_{dL}} = 2.3^{\pm 1} \times 10^{-3}, \\
\boldsymbol{\delta g^Z_{uR}} = -3.6^{\pm 3.5} \times 10^{-3}, \quad \boldsymbol{\delta g^Z_{dR}} = 16.0^{\pm 5.2} \times 10^{-3},
\end{aligned}$$

[d] This is not necessary true for effects that grow with energy and can be measured in VH associated production processes, as discussed in Ref. [14].

with a correlation matrix reported in Ref.[15]; from LEP2 data, on the other hand, we obtain[e] $\delta g_{1,Z} = -0.05^{+0.05}_{-0.07}$ and $\delta\kappa_\gamma = 0.05^{+0.04}_{-0.04}$.

On the other hand, the following effects, which also affect Higgs and EW physics, do not interfere with the SM:

$$\Delta\mathcal{L}_R^W = \boldsymbol{\delta g_R^W} \frac{\hat{h}^2}{v^2} W_\mu^+ \bar{u}_R \gamma^\mu d_R + \text{h.c.}, \tag{12}$$

$$\begin{aligned}\Delta\mathcal{L}_{\text{dipole}}^V = & \frac{eY_q\hat{h}}{m_W^2}\Big[\boldsymbol{\delta\kappa_q^G}\frac{g_s}{e}\bar{q}_L T^A\sigma^{\mu\nu}q_R G^A_{\mu\nu} + \boldsymbol{\delta\kappa_q^A}(T_3\bar{q}_L\sigma^{\mu\nu}q_R A_{\mu\nu} + \frac{s_{\theta_W}}{\sqrt{2}}\bar{u}_L\sigma^{\mu\nu}d_R W^+_{\mu\nu}) \\ & + \boldsymbol{\delta\kappa_q^Z}(T_3\bar{q}_L\sigma^{\mu\nu}q_R Z_{\mu\nu} + \frac{c_{\theta_W}}{\sqrt{2}}\bar{u}_L\sigma^{\mu\nu}d_R W^+_{\mu\nu}) + \text{h.c.}\Big],\end{aligned} \tag{13}$$

for quarks $q = u, d$, where the coefficients are assumed to be real and T_3 denotes weak isospin (and similarly for leptons). Here $\boldsymbol{\delta g_R^W}$ is expected to be suppressed by both the down- and up-type Yukawas in a MFV expansion so that (together with the fact that it doesn't interfere with the SM and its contribution is therefore suppressed in inclusive quantities) it can be neglected. On the other hand the $\boldsymbol{\delta\kappa_q^V}$ can be measured in dipole-type experiments and we omit the result here.

Finally $\mathcal{L}_6$ includes interactions that do not involve the Higgs field. In particular

$$\Delta\mathcal{L}_{\lambda_\gamma} = \frac{i\boldsymbol{\lambda_\gamma}}{m_W^2}\left[(eA^{\mu\nu} + gc_{\theta_W}Z^{\mu\nu})W^{-\rho}_\nu W^+_{\rho\mu}\right], \quad \Delta\mathcal{L}_{3G} = \frac{\boldsymbol{\kappa_{3G}}}{m_W^2} g_s \epsilon_{ABC} G^{A\,\nu}_\mu G^B_{\nu\rho} G^{C\,\rho\mu}. \tag{14}$$

and four-fermion interactions, which can be found in [10]. LEP2 data[15] gives $\lambda_\gamma = 0.00^{+0.07}_{-0.07}$, while $\boldsymbol{\kappa_{3G}}$ and four-fermion interactions involving quarks can be constrained using dijet searches at the LHC [17]. Interactions involving leptons and quarks can be constrained at LEP [18] and LHC[6,19].

In summary, Eqs. (1-6) together with Eqs. (7-10), Eqs. (12,13), Eq. (14) and the four-fermion interactions, offer a complete parametrization of all BSM effects accessible at the leading order in an expansion in inverse powers of the new physics scale. They are organized in such a way that experimental (*input*) constraints can be readily implemented and the physical consequences quickly extrapolated, as we show in the next section.

3 Consequences

The main predictions from this analysis are the following.[6,5] First of all, from Eqs. (7,8) it is clear that the Wff and Zff vertices are related at the level of $\mathcal{L}_6$, while the W dipole-type interaction for the fermions are related to those of A and Z as can be read from Eq. (13). Furthermore, there are only 3 types of CP-conserving TGC, characterized by[20] δg_1^Z, $\delta\kappa_\gamma$ and λ_γ, while QGC are related to them through

$$\boldsymbol{\delta g_1^Z} = \frac{\delta g^Z}{g^Z_{SM}} = \frac{\delta g^{WW}}{2c^2_{\theta_W}g^{WW}_{SM}} = \frac{\delta g^{ZZ}}{2g^{ZZ}_{SM}} = \frac{\delta g^{\gamma Z}}{g^{\gamma Z}_{SM}} \tag{15}$$

and Eq. (14). Finally, there are only 8 Higgs BSM primary effects (for one family), given in Eqs. (1-6), while all other Higgs interactions can be written as function of the parameters of $\mathcal{L}_6$ discussed so far. An interesting example is the differential distribution of $h \to Vff$[21,22,23,24,25,26,27,28,29], whose amplitude is generically written as

$$\mathcal{M}(h \to VJ_f) = (\sqrt{2}G_F)^{1/2}\epsilon^{*\mu}(q)\,J_f^{V\nu}(p)\left[A_f^V\,\eta_{\mu\nu} + B_f^V\,(p\cdot q\,\eta_{\mu\nu} - q_\mu p_\nu)\right], \tag{16}$$

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

$$A_f^V = a_f^V + \widehat{a}_f^V\frac{p^2+M_V^2}{p^2-M_V^2}, \qquad B_f^V = b_f^V\frac{1}{p^2-M_V^2} + \widehat{b}_f^V\frac{1}{p^2} \qquad (\widehat{b}_f^V = 0 \text{ for } V = W). \tag{17}$$

[e] The extraction of these parameters from data, is complicated by the limited experimental information available, as discussed in Refs. [15,16].

Now, the coefficients a_f^V and b_f^V are associated with Lagrangian structures, such as $hV_{\mu\nu}V^{\mu\nu}$, whose coefficient in $\mathcal{L}_6$ can be readily read from the expressions in the previous section. For the case of $h \to Z\bar{l}l$, we find

$$\frac{\delta a^Z_{l_L}}{a^Z_{l_L}} \in [-0.2, 0.1]\,, \qquad \frac{\delta \widehat{a}^Z_{l_L}}{\widehat{a}^Z_{l_L}} \in [-8,7]\times 10^{-2}\,,\ \ b^Z_{l_L} \in [-2,5]\times 10^{-2}, \qquad \widehat{b}^Z_{l_L} \in [-2,5]\ ,\times 10^{-2},$$
$$\frac{\delta a^Z_{l_R}}{a^Z_{l_R}} \in [-0.2, 0.3]\,, \qquad \frac{\delta \widehat{a}^Z_{l_R}}{\widehat{a}^Z_{l_R}} \in [-8,7]\times 10^{-2},\ \ b^Z_{l_R} \in [-3,2]\times 10^{-2}, \qquad \widehat{b}^Z_{l_R} \in [-2,5]\times 10^{-2}\,.$$

Although the allowed range in $a^Z_{l_{L,R}}$ is quite large, we notice that their impact on the total amplitude, when summed over lepton chiralities, is much smaller, $2\sum_{l=l_L,l_R} g^l_Z a^Z_l / \sum_{l=l_L,l_R}(g^l_Z)^2 \in [-6,4]\times 10^{-2}$. This implies that the expected BSM modification in the differential distribution of Higgs decay is already fairly constrained: our analysis sets the goal for future Higgs physics experiments to be competitive.

Interestingly, this differential distribution, although not directly tested by experimental collaborations so far, has been probed by measurements of the custodial parameter λ_{WZ} during LHC Run1. In fact, momentum-dependent deformations of the $hV\bar{f}f$ coupling behave differently when tested in $h \to W\bar{f}f'$ or $h \to Z\bar{f}f$, because of the difference between m_Z and m_W. In some sense, the custodial parameter λ_{WZ} is sensitive to the SM custodial symmetry breaking, through custodial-preserving momentum-dependent interactions. Through our analysis we find

$$\lambda_{WZ} - 1 \ \equiv\ \frac{\Gamma(h\to WW)}{\Gamma^{\rm SM}(h\to WW)}\frac{\Gamma^{\rm SM}(h\to ZZ)}{\Gamma(h\to ZZ)} - 1 \simeq 0.8 g_1^Z - 0.1\kappa_\gamma - 1.6\kappa_{Z\gamma} \in [-5,6]\times 10^{-2}\,,$$

We see that Eq. (18) puts a bound on λ_{WZ} stronger than the experimental limit [8]: $(\lambda_{WZ}-1) \in [-0.45, 0.15]$.

4 Conclusions

The SM EFT motivates SM precision tests, providing a framework in which searches for departures from the SM can be interpreted as searches for new physics and can be compared with direct searches of explicit models. In a bottom-up approach, the parameters characterizing the leading BSM piece of this effective Lagrangian, $\mathcal{L}_6$, can be fixed through the most precise SM precision tests. Then, since the parameters in $\mathcal{L}_6$ are less than the observables that are modified by $\mathcal{L}_6$, we can relate different observables and extract concrete, but generic, predictions. This task is facilitated by writing $\mathcal{L}_6$ in the *BSM Primaries* basis, where observables can be written in terms of other observables. Using this procedure, we have provided a quantitative prediction for the expected variation of the differential distribution of $h \to V\bar{f}f$ decays, for the custodial parameter λ_{WZ}, for the W couplings to fermions, for quartic gauge couplings and for dipole-type interactions involving the W-boson. These relations can be used to understand which observables deserve more attention in future experiments and which, instead, are already well measured.

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