Signatures of a nonstandard Higgs from flavor physics

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We examine the constraints coming from incorporating the full Standard Model gauge symmetry into the effective field theory description of flavor processes, using semileptonic decays as paradigmatic examples. Depending on the dynamics triggering electroweak symmetry breaking, different patterns of correlations between the Wilson coefficients arise. Interestingly, this implies that flavor experiments are capable of shedding light upon the nature of the Higgs boson without actually requiring Higgs final states. Furthermore, the observed correlations can simplify model-independent analyses of these decays.

I. Introduction. The discovery of the Higgs boson at the LHC \cite{HiggsDiscovery} and the current knowledge of its properties have once more confirmed that the Standard Model (SM) is an excellent low-energy description of the electroweak interactions. The precise nature of the Higgs boson – which is linked to the characteristics of potential new physics (NP) – is however an issue that remains to be settled; this is one of the main goals of Run-II at the LHC, using the analysis of Higgs couplings in multi-Higgs production processes. However, recent studies \cite{Choudalakis:2013vwa, Choudalakis:2014yua} draw rather pessimistic conclusions regarding the capability of the LHC to discriminate, e.g., weakly- from strongly-coupled NP scenarios from double Higgs production.

In this paper we argue that flavor processes have the potential to test the dynamics triggering electroweak symmetry breaking (EWSB), despite being unable to probe Higgs couplings directly. In that sense, flavor experiments can provide valuable information on the Higgs and should be seen as complementary to the LHC effort to identify its properties. Turning the argument around, interesting consequences arise for flavor physics when assuming weakly-coupled NP.

For illustration we will consider heavy-flavor semileptonic processes, specifically the flavor-changing neutral current (FCNC) processes $D \to D' \ell^+ \ell^-$ and $U \to U' \ell^+ \ell^-$ as well as the charged-current processes $D \to D \ell^+ \ell^-$ (effectively including $U \to D \ell^+$), where $D^{(1)}$ and $U^{(1)}$ are down-type and up-type quarks, respectively, and $\ell = e, \mu, \tau$. For definiteness we will focus on $b \to s \ell \ell$, $c \to u \ell \ell$ and $b \to c \tau \nu$ as representatives of the different groups. At hadronic scales $\Lambda \ll M_\text{H}$ the effective field theory for down-quark semileptonic processes reads \cite{Buchalla:1994ui, Buchalla:1995vs, Buchalla:1995vs}.

\begin{equation}
\mathcal{L}^{b\to s\ell\ell}_{\text{eft}} = \frac{4G_F}{\sqrt{2}} \lambda_{bs} \frac{e^2}{(4\pi)^2} \sum_{i}^{12} C_i \mathcal{O}_i,
\end{equation}

where $\lambda_{bs} = V_{ub}V_{us}^\ast$ and the operators are defined as

\begin{equation}
\mathcal{O}_i^{(1)} = \frac{m_b}{c} (\bar{s} \sigma^{\mu\nu} P_{R(L)} b) F_{\mu\nu},
\end{equation}

\begin{equation}
\mathcal{O}_i^{(2)} = (\bar{s} \gamma_\mu P_{L(R)} b) \bar{\ell} \gamma^\mu \ell,
\end{equation}

\begin{equation}
\mathcal{O}_i^{(3)} = (\bar{s} P_{R(L)} b) \bar{\ell} \gamma_3 \ell,
\end{equation}

\begin{equation}
\mathcal{O}_i^{(5)} = (\bar{s} \sigma_{\mu\nu} b) \bar{l} \sigma^{\mu\nu} l,
\end{equation}

For up-quark transitions, one likewise finds

\begin{equation}
\mathcal{L}^{c\to u\ell\ell}_{\text{eft}} = \frac{4G_F}{\sqrt{2}} \lambda_{bu} \frac{e^2}{(4\pi)^2} \sum_{i}^{12} C_i \mathcal{O}_i^{(a)}(a),
\end{equation}

where $\lambda_{bu} = V_{cb}V_{ub}^\ast$ and the operators are obtained by the trivial flavor replacements $(b; s) \to (c; u)$ in Eq. \textsuperscript{(2)}.

Finally, charged-current decays are described by

\begin{equation}
\mathcal{L}^{b\to c\ell\nu}_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} V_{cb} \sum_{j} c_{j} \mathcal{O}_j,
\end{equation}

where the corresponding operators are defined as follows:

\begin{equation}
\mathcal{O}_{V_{L,R}} = (\bar{c} \gamma^\mu P_{L(R)} b) \bar{\tau} \gamma_\mu \nu,
\end{equation}

\begin{equation}
\mathcal{O}_{S_{L,R}} = (\bar{c} P_{L(R)} b) \bar{\tau} \nu,
\end{equation}

\begin{equation}
\mathcal{O}_{T} = (\bar{c} \sigma^{\mu\nu} P_{L} b) \bar{\tau} \sigma_{\mu\nu}\nu.
\end{equation}

The structure of the occurring operators is constrained by the strong and electromagnetic symmetries. Information about the full electroweak symmetry can be implemented by matching the previous set of operators to effective field theories (EFTs) valid at the electroweak scale (see Refs. \cite{Buchalla:1995vs, Buchalla:1995vs} for earlier work in this direction). However, the specific form of the EFT at the electroweak scale strongly depends on the nature of the Higgs boson. For a conventional SM Higgs, the scalar sector is described by a linear sigma model, where the Higgs and the electroweak Goldstone bosons belong to a weak doublet. In order to test the SM Higgs hypothesis, one needs a more general framework where the Goldstone bosons can be decorated from the Higgs particle. Such a framework is provided by a nonlinear representation, see e.g. Ref. \textsuperscript{8} for a general discussion. In its minimal implementation,
the Higgs is a scalar singlet and the electroweak Goldstone modes are contained in a matrix $U$ transforming as a bifundamental of $SU(2)_L \times SU(2)_R$, $U \rightarrow g_U U g_U^T$. The linear and nonlinear representations correspond to markedly different dynamical pictures of the mechanism triggering EWSB. The linear EFT is suitable to describe weakly-coupled extensions of the SM, yielding an expansion in canonical dimensions. The nonlinear EFT is in contrast aimed at strongly-coupled dynamics, and requires a loop expansion (or equivalently an expansion in chiral dimensions).

Matching the nonlinear EFT to the flavor basis will therefore serve a two-fold purpose: (i) it will ensure a model-independent implementation of the electroweak symmetry for flavor processes and (ii) it will provide a framework to characterize departures from an SM Higgs.

2. Matching at the electroweak scale. We match the NLO nonlinear operator basis at the electroweak scale into the flavor EFT as a bifundamental of $SU(2)_L \times SU(2)_R$, $U \rightarrow g_U U g_U^T$. The situation is substantially different for the scalar and tensor sectors. For the down-quark sector one finds

$$\delta C_{7(d)}^{(t)} = \frac{8\pi^2}{m_t \lambda_{ts}} \frac{v^2}{\Lambda^2} \left[ c_X^2 + c_X^4 \delta \right],$$

$$\delta C_{7(u)}^{(t)} = \frac{8\pi^2}{m_u \lambda_{bu}} \frac{v^2}{\Lambda^2} \left[ c_X^2 + c_X^4 \delta \right],$$

while for the vector sector one finds (defining $N_{NC}^{(d)} = 4\pi^2 \lambda_{ts} \frac{v^2}{\Lambda^2}$ and $N_{NC}^{(u)} = 4\pi^2 \lambda_{bu} \frac{v^2}{\Lambda^2}$)

$$\delta C_{9,10}^{(q)} = \delta \left[ \left( \frac{c(q)}{C_{LL}} + C_{LL} \right) \pm 4 g_{V,A} \frac{\Lambda^2}{v^2} C_{V,L}^{(q)} \right],$$

$$\delta C_{9,10}^{(q)} = \delta \left[ \left( \frac{c(q)}{C_{RR}} + C_{RR} \right) \pm 4 g_{V,A} \frac{\Lambda^2}{v^2} C_{V,R}^{(q)} \right].$$

The main conclusion to be drawn, independently of the EWSB mechanism, is that for the dipole operators and the vector sector invariance under the electroweak symmetry does not add information compared to only imposing electromagnetic and strong invariance, i.e., the number of independent operators does not get reduced. This is not unexpected, since there are strong indications that these sectors are effectively decoupled from the mechanism of EWSB. In fact, in the nonlinear framework they only appear as finite counterterms, i.e., they are not needed to renormalize the EFT. As a result they are expected to be rather insensitive to electroweak physics in general and to the existence of the Higgs in particular.

The situation is substantially different for the scalar and tensor sectors. For the down-quark sector one finds

$$C_{S,P}^{(d)} = N_{NC}^{(d)} \left[ c_{S,P}^{(d)} + \delta c_{V}^{(d)} \right],$$

$$C_{T}^{(d)} = N_{NC}^{(d)} \left[ c_{T}^{(d)} + \delta c_{T}^{(d)} \right],$$

where $c_{S,P}^{(d)} = 2(c_{LR}^{(d)} - c_{LR}^{(d)})$. Notice that in the linear case, i.e., assuming a standard Higgs, one finds

$$C_{S}^{(d)} = -C_{P}^{(d)} , \quad C_{S}^{(d)} = C_{P}^{(d)} , \quad C_{T}^{(d)} = C_{T}^{(d)} = 0 ,$$

as already noted in [5]. Deviations from these relations are expected to arise at the percent level, through Higgs-exchange diagrams (which are NLO but numerically suppressed by a small Yukawa coupling) and also by $O(v^2/\Lambda^2)$ effects induced by NNLO operators, all of which can be safely neglected. However, it is important to note that they are not predictions of electroweak symmetry alone. For a nonstandard Higgs, one finds that the pattern of correlations is already broken at NLO: the scalar sector is fully decorrelated and $C_{T}$ and $C_{T5}$ no longer vanish. These decorrelations are caused by the operators $O_{Y_{ij}}$ which are, contrary to the remaining operators listed in the Appendix, characterized by having a fermion content with a nonvanishing total hypercharge.

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3 Here and in the following, primed coefficients are related to unprimed coefficients by taking the same sample operator listed in the Appendix, reverting the quark flavor indices and applying hermitean conjugation.

4 If NP lies at the TeV scale, there is no generic reason to expect it to be coupled to the SM fermions with Yukawa-like patterns. The local contributions therefore dominate over the Higgs-exchange diagrams in general.
this is of course compensated by the hypercharge of the $U$ field. For a standard Higgs such operators can only appear at NNLO and therefore provide a clean way of fingerprinting the nature of the Higgs boson.

A similar situation is encountered in the up-quark sector, where we obtain

\[
\begin{align*}
C_{S,P}^{(u)} &= N_{\text{NC}}^{(u)} \left[ \hat{c}_S^{(u)} + \hat{c}_Y \right], \\
C_{T}^{(u)} &= N_{\text{NC}}^{(u)} \left[ \hat{c}_T^{(u)} + \hat{c}_T \right], \\
C_{T5}^{(u)} &= N_{\text{NC}}^{(u)} \left[ \hat{c}_T^{(u)} - \hat{c}_T \right], \\
C_{S,T}^{(u)} &= N_{\text{NC}}^{(u)} \left[ \hat{c}_S^{(u)} + \hat{c}_Y \right], \\
C_{S,P}^{(u)} &= N_{\text{NC}}^{(u)} \left[ \hat{c}_S^{(u)} - \hat{c}_T \right],
\end{align*}
\]

(12)

with $c_r^{(u)} = -c_{S1}^{(u)} + c_{S4}^{(u)}$ and $c_t^{(u)} = -c_{S2}^{(u)} + c_{S4}^{(u)}$. A standard Higgs would in this case predict

\[
C_{S,P}^{(u)} = C_{P}^{(u)}, \quad C_{S,P}^{(u)} = -C_{P}^{(u)}. \quad (13)
\]

Instead, for a nonstandard Higgs, Eqs. (13) are no longer satisfied due to the presence of $\hat{c}_Y$. Due to the hypercharge structure of the transition, contributions to the tensor sector appear already at NLO in both the linear and nonlinear EFTs and cannot be used to discriminate between the two.

We now turn our attention to charged-current processes. For $b \rightarrow c\ell\nu$ transitions the matching between the flavor and electroweak EFTs reads ($N_{\text{CC}} = \frac{1}{2\alpha_{\text{em}}^2}$)

\[
\begin{align*}
C_{V_L} &= -N_{\text{CC}} \left[ C_L + \frac{2}{v^2} C_V + \frac{5}{2\alpha_{\text{em}}^2} C_{V7} \right], \\
C_{V_R} &= -N_{\text{CC}} \left[ C_L + \frac{2}{v^2} C_V + \frac{5}{2\alpha_{\text{em}}^2} C_{V7} \right], \\
C_{S_L} &= N_{\text{CC}} \left[ c_{S1}^{(u)} + c_{S4}^{(d)} \right], \\
C_{S_R} &= 2N_{\text{CC}} \left( c_{LR4} + c_{LR8} \right), \\
C_{T} &= -N_{\text{CC}} \left[ c_{S2}^{(u)} + c_{S6}^{(u)} \right],
\end{align*}
\]

(14)

where $C_L = 2c_{LL2} - \hat{c}_{LL6} + \hat{c}_{LL7}$ and $\hat{C}_R = -\frac{1}{2} \hat{c}_Y$. To simplify the notation, flavor rotation matrices have been absorbed into the NLO coefficients of the EFT. This is always possible when considering a single transition. However, when relating different processes, relative flavor rotations have to be taken into account explicitly, see below.

While in this case for both electroweak EFTs the full basis is reproduced with independent coefficients, correlations between different processes appear in the linear case with interesting consequences: (i) the absence of a direct four-fermion operator contribution to $C_{V_L}$ implies lepton-flavor universality in that sector (inherited from the $W$-to-fermion couplings in the SM), as already noted in [3]; (ii) the scalar sectors of the charged and neutral processes are related, see e.g. [15]. For instance,

\[
\sum_{U=u,c,t} \lambda_{Us} C_{S,P}^{(u)} = -\frac{e^2}{8\pi^2} \lambda_3 C_S^{(d)} \quad (15)
\]

between $b \rightarrow U\ell\nu$ and $b \rightarrow s\ell\ell$, where we have taken into account the relative flavor rotation to the quark mass eigenstates. Similar relations hold for all processes related by $SU(2)_L$. We remark that this is non-trivial for the coefficients in the flavor EFT which generally receive contributions from several electroweak invariant operators. Incidentally, we note that Eq. (15) implies that bounds from FCNC processes severely constrain the size of new physics in their $SU(2)$-related charged-current processes.

We stress again that the correlations (i) and (ii) appear exclusively in the standard Higgs scenario and therefore can be used to test the nature of the Higgs boson from flavor physics.

3. Nontrivial hypercharge operators from TeV physics. The operators that most prominently distinguish weakly-from strongly-coupled Higgs scenarios are $\hat{O}_{Y_j}$, i.e. four-fermion structures with nonvanishing fermionic hypercharge. In a weakly-coupled scenario such structures can only appear at NNLO, e.g. $(q\bar{q}d)(\bar{l}_c\ell e)$ or $(\bar{l}_c\ell d)(\bar{u}_b\bar{q}1')$, while in the nonlinear case they are present already at NLO. Such operators can be generated in simple models of heavy scalar exchanges, as we shall illustrate here.

As an example we consider an extension of the SM where EWSB is driven by a strongly-coupled sector with the addition of two TeV-scale states: a scalar $\phi$, singlet under the SM group, and a colored scalar $\Phi$, transforming as $(3,2)$–$2/3$. These heavy states could be fundamental fields or composite states of the strong sector. Since we are interested in four-fermion operators, we will concentrate on their couplings to fermionic scalar currents:

\[
\begin{align*}
L_{\text{int}}(\phi, \Phi) &= \lambda_3 \frac{\bar{q}U P_+ \phi \sigma + \lambda_4 \bar{q}U P_- \phi \tau + \lambda_5 \bar{q}U P_- \phi \eta}{\lambda_3 \frac{\bar{q}U P_+ \Phi r + \lambda_6 \bar{q}U P_- \Phi r + \lambda_7 \bar{q}U P_- \Phi q}{h.c.}. \quad (16)
\end{align*}
\]

For simplicity the coefficients are assumed to be real. Integrating out the heavy scalars, the resulting EFT includes the terms

\[
\begin{align*}
\mathcal{L}_{\text{eff}} &= \sum_{q} \frac{\lambda_3 \lambda_4}{2m_\phi^2} \hat{O}_{Y3} + \frac{\lambda_3 \lambda_5}{2m_\phi^2} \hat{O}_{S3} + \frac{\lambda_4 \lambda_5}{2m_\phi^2} \hat{O}_{Y1} \\
&- \frac{\lambda_3 \lambda_7}{8m_\phi^2} (\hat{O}_{LR4} - \hat{O}_{LR5}) - \frac{\lambda_4^2}{8m_\phi^2} (\hat{O}_{LR2} - \hat{O}_{LR6}) \\
&- \frac{\lambda_7^2}{8m_\phi^2} (\hat{O}_{LR3} - \hat{O}_{LR7}) - \frac{\lambda_3^2}{8m_\phi^2} (\hat{O}_{LR1} - \hat{O}_{LR5}) \\
&+ \frac{\lambda_3 \lambda_7}{2m_\phi^2} \hat{O}_{Y4} - \frac{\lambda_4 \lambda_5}{4m_\phi^2} \left( \hat{O}_{S3} - \frac{1}{4} \hat{O}_{S4} \right) \\
&- \frac{\lambda_3 \lambda_7}{4m_\phi^2} \left( \hat{O}_{Y1} - \frac{1}{4} \hat{O}_{Y2} \right) + h.c., \quad (17)
\end{align*}
\]

therefore generating explicitly the operators that violate the relations [11], [13] and [15].

4. Conclusions. We have analyzed the impact of the full electroweak symmetry at hadronic scales, extending
the work done in Refs. [6, 7] and including nonstandard Higgs scenarios. We have shown that correlations between different coefficients are linked to assuming an SM (i.e., weak doublet) Higgs field. Their violation would point at a nonstandard Higgs, presumably of a composite nature. Flavor processes therefore provide valuable information on the dynamics responsible for EWSB. Furthermore, we have explored the consequences of assuming a linearly realized electroweak symmetry for Higgs scenarios. We have shown that correlations between the corresponding model-independent analyses. Specifically, in addition to Eqs. (11) for $D \to D'\ell\ell$ decays, already reported in Ref. [7], we have found Eqs. (13) for $U \to U'\ell\ell$ decays. In charged-current decays, the electroweak symmetry implies lepton-flavor universality for the right-handed vector currents. Furthermore, we pointed out the phenomenological significance of SU(2) relations between FCNC and charged-current semileptonic decays, e.g. Eq. (14).

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Appendix. In the following we list the subset of relevant operators for semileptonic processes, extracted from the NLO operator basis worked out in Ref. [10]. Notational changes have been introduced to ease the comparison with the linear basis in Ref. [12]. For simplicity, we will work in unitary gauge, where $U = 1$. Below we will however keep $U$ in order to make the invariance under the SM gauge symmetry transparent. $\tau_3 = U\tau_3 U^\dagger$ and $\tau_\pm = U\left(\tau_1 \pm i\tau_2\right) U^\dagger$ are chirally-dressed Pauli matrices, and $L_\mu \equiv uD_j U^\dagger$.

The subset of operators relevant for the electromagnetic dipole operators are

$$\mathcal{O}_{X_{1,2}} = g' \sigma^{\mu\nu} U P_{\nu} R_{\mu}, \quad \mathcal{O}_{X_{3,4}} = g \sigma^{\mu\nu} U P_{\nu} r(\tilde{\tau}_3 W_{\mu}),$$

while the one for the vector operators reads

$$\mathcal{O}_{V_{1,2}} = \tilde{\gamma}^{\mu} q \tilde{v}_{\mu} l, \quad \mathcal{O}_{V_{3}} = \tilde{\gamma}^{\mu} \tilde{\tau}_3 q \tilde{\tau}_3 l, \quad \mathcal{O}_{V_{4}} = \tilde{\gamma}^{\mu} d(\tilde{\tau}_3 L_\mu), \quad \mathcal{O}_{V_{5}} = \tilde{\gamma}^{\mu} \tilde{\tau}_3 q \tilde{\tau}_3 L_\mu, \quad \mathcal{O}_{V_{6}} = \tilde{\gamma}^{\mu} d(\tilde{\tau}_3 L_\mu),$$

Finally, operators relevant for the scalar and tensor sectors at the electroweak scale are

$$\mathcal{O}_{S_{1,2}} = \epsilon_{ij} \tilde{\tau}_i \epsilon^{ij} q, \quad \mathcal{O}_{T_{1,2}} = \tilde{\gamma}^{\mu} \tilde{\tau}_3 q \tilde{\tau}_3 l, \quad \mathcal{O}_{T_{3}} = \tilde{\gamma}^{\mu} \tilde{\tau}_3 d \tilde{\tau}_3 l, \quad \mathcal{O}_{T_{4}} = \tilde{\gamma}^{\mu} \tilde{\tau}_3 d \tilde{\tau}_3 q,$$

Flavor family indices have been omitted.

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