Effective field theory with two Higgs doublets

Andreas Crivellin,[a] Margherita Ghezzi[a] and Massimiliano Procura[b,1]

[a] *Paul Scherrer Institut,*
 CH-5232 Villigen PSI, Switzerland
[b] *Theoretical Physics Department, CERN,*
 Geneva, Switzerland

 E-mail: andreas.crivellin@cern.ch, margherita.ghezzi@psi.ch,
 massimiliano.procura@cern.ch

ABSTRACT: In this article we extend the effective field theory framework describing new physics effects to the case where the underlying low-energy theory is a Two-Higgs-Doublet model. We derive a complete set of independent operators up to dimension six assuming a Z_2-invariant CP-conserving Higgs potential. The effects on Higgs and gauge boson masses, mixing angles in the Higgs sector as well as couplings to fermions and gauge bosons are computed. At variance with the case of a single Higgs doublet, we find that pair production of SM-like Higgses, arising through dimension-six operators, is not fixed by fermion-fermion-Higgs couplings and can therefore be sizable.

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[1]On leave from the University of Vienna, Austria.

Contents

1 Introduction

Adding a second SU(2) doublet scalar [1] represents one of the simplest possible extensions of the Standard Model (SM) of strong and electroweak (EW) interactions. Two-Higgs-Doublet models (2HDMs) have been under extensive investigation for a long time (see for example ref. [2] for an introduction and ref. [3] for a more recent review article). There are several reasons for this interest. Firstly, these models have a rather small number of free parameters, which makes phenomenological analyses quite predictive. Additional motivation for 2HDMs is provided by axion models [4], where a global U(1) Peccei-Quinn symmetry is introduced to eliminate a CP-violating term in the QCD Lagrangian [5]. Such a symmetry is only possible in scenarios with at least two Higgs doublets. Furthermore, the amount of CP violation obtained through the introduction of a second Higgs doublet can be large enough to account for the baryon asymmetry of the universe [6]. 2HDMs can also explain the anomalies observed in tauonic B decays [7–9]. Finally, strong motivation for studying 2HDMs is provided by the Minimal Supersymmetric Standard Model (MSSM) where supersymmetry enforces the introduction of a second Higgs doublet [10–12] due to the holomorphicity of the superpotential.

According to present collider bounds, the additional Higgs bosons contained in these models are still allowed to have masses around the EW scale [13, 14].[1] Therefore, in an effective field theory (EFT) approach, it is natural to consider the additional Higgs doublet as a dynamical degree of freedom like the SM fields. In this article, we shall consider 2HDMs as EFTs valid up to a high-energy scale $\Lambda \gg M_H$ where additional dynamical degrees of freedom enter. The MSSM with heavy SUSY partners but light Higgs doublets is one

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[1]An exception is the 2HDM of type II where the bound from $b \to s\gamma$ forces the charged Higgs mass (which only differs from the other Higgs masses by terms of order v^2) to be larger than 400 GeV [15].

example of such a theory, but also the $L_\mu - L_\tau$ model of refs. [16, 17] reduces to a 2HDM if the Z' and the $L_\mu - L_\tau$-breaking singlet are heavy.

In general, any theory of new physics (NP) super-seeding the 2HDM at higher energies must satisfy the following requirements (in close analogy to the SM case):

(i) Its gauge group contains the SM gauge group $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ as a subgroup.

(ii) It contains two Higgs doublets as dynamical degrees of freedom, either as fundamental or composite fields.

(iii) At low energies it reproduces the 2HDM, barring the existence of weakly coupled *light* particles, like axions or sterile neutrinos (in which case the EFT should include in addition these particles as dynamical degrees of freedom).

In this approach, heavier NP particles are integrated out and their effects are parameterized in terms of Wilson coefficients of higher-dimension operators suppressed by inverse powers of Λ. In our 2HDM case we have

$$\mathcal{L}_{2HDM} = \mathcal{L}_{2HDM}^{(4)} + \frac{1}{\Lambda} \sum_k C_k^{(5)} Q_k^{(5)} + \frac{1}{\Lambda^2} \sum_k C_k^{(6)} Q_k^{(6)} + \mathcal{O}\left(\frac{1}{\Lambda^3}\right). \tag{1.1}$$

Here $\mathcal{L}_{2HDM}^{(4)}$ is the standard renormalizable 2HDM Lagrangian (to be specified in the next section) which contains only dimension-two and dimension-four operators. $Q_k^{(5)}$ generalize the Weinberg operator [18] giving rise to neutrino masses and $Q_k^{(6)}$ denote the dimension-six operators. $C_k^{(5)}$ and $C_k^{(6)}$ are their dimensionless Wilson coefficients. In this paper we neglect the effects of operators of dimension seven and higher, which are suppressed by at least three powers of Λ.

The EFT approach to parameterize NP effects through higher-dimension operators built with SM fields (SM-EFT) has been used for a long time [19–22]. More recently, a complete and minimal basis of dimension-six effective operators in the SM has been established [23] and the SM-EFT has received a lot of interest concerning its phenomenological applications (see for example [24–28]), mainly in the context of Higgs physics [29–37] but also flavor physics [38–47]. The possibility to extract constraints on the Wilson coefficients both from Higgs measurements at the LHC and from electroweak precision observables has been exploited [48–55] and an effort is currently made to perform the SM-EFT analysis at next-to-leading order in perturbation theory [56–67].

The purpose of this paper is to extend the SM-EFT approach to the case of two Higgs doublets. 2HDMs are considered here as low-energy theories and using our framework the effects of heavier NP particles can be clarified and studied in a systematic way. For example, effective operators in 2HDMs have been recently considered to account for a diphoton excess in LHC data [68, 69] with the outcome that an explanation requires to extend the field content of pure 2HDMs [70–79].[2]

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[2]Prior to the LHC diphoton excess, effective operators for two-photon processes in 2HDMs were already discussed in ref. [80].

	fermions					scalars
field	l_{Lp}^j	e_{Rp}	$q_{Lp}^{\alpha j}$	u_{Rp}^{α}	d_{Rp}^{α}	φ_1^j, φ_2^j
hypercharge Y	$-\dfrac{1}{2}$	-1	$\dfrac{1}{6}$	$\dfrac{2}{3}$	$-\dfrac{1}{3}$	$\dfrac{1}{2}$

Table 1. The matter content of the 2HDM. l (e) is the lepton doublet (singlet), u and d the right-handed up and down quark singlets and q the quark doublet. Here, $j = 1, 2$, $\alpha = 1, 2, 3$, and $p = 1, 2, 3$ stand for isospin, color and generation indices, respectively.

This article is organized as follows. In section 2 we will introduce our notation and conventions. The complete list of operators up to dimension six before spontaneous EW symmetry breaking will be given in section 3. We will then discuss EW symmetry breaking and the definition of the physical basis with diagonal mass matrices in section 4. In section 5 we will conclude and point out an interesting phenomenological application of our formalism.

2   Notation and conventions

In this section we establish our notation and conventions following ref. [23] and ref. [3]. The renormalizable 2HDM Lagrangian before spontaneous EW symmetry breaking, reads

$$
\begin{aligned}
\mathcal{L}_{2HDM}^{(4)} = {}& -\frac{1}{4} G_{\mu\nu}^A G^{A\mu\nu} - \frac{1}{4} W_{\mu\nu}^I W^{I\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} \\
& + (D_\mu \varphi_1)^\dagger (D^\mu \varphi_1) + (D_\mu \varphi_2)^\dagger (D^\mu \varphi_2) \\
& - V(\varphi_1, \varphi_2) + i \left(\bar{l} \slashed{D} l + \bar{q} \slashed{D} q + \bar{u} \slashed{D} u + \bar{d} \slashed{D} d \right) + \mathcal{L}_Y \,,
\end{aligned}
\tag{2.1}
$$

where φ_1 and φ_2 are the two Higgs doublets, $A = 1 \ldots 8$ ($I = 1 \ldots 3$) labels the SU(3) (SU(2)) gauge bosons while $B_{\mu\nu}$ is the hypercharge field strength tensor. Note that we omitted the QCD θ-term here. The fermion fields and their charges and representations are shown in table 1. For the sake of simplicity, chirality indices are suppressed in the following. The conventions for covariant derivatives are fixed e.g. by

$$
(D_\mu q)_{\alpha j} = \left(\partial_\mu + i g_s T_{\alpha\beta}^A G_\mu^A + i g S_{jk}^I W_\mu^I + i g' Y_q B_\mu \right) q^{\beta k} \,.
\tag{2.2}
$$

with $T^A = \frac{1}{2} \lambda^A$ and $S^I = \frac{1}{2} \tau^I$ denoting the SU(3) and SU(2) generators, and λ^A (τ^I) the Gell-Mann (Pauli) matrices. It is useful to define the following Hermitian derivative terms:

$$
\varphi^\dagger i \overleftrightarrow{D}_\mu \varphi \equiv i \varphi^\dagger D_\mu \varphi - i (D_\mu \varphi)^\dagger \varphi \qquad \text{and} \qquad \varphi^\dagger i \overleftrightarrow{D}_\mu^I \varphi \equiv i \varphi^\dagger \tau^I D_\mu \varphi - i (D_\mu \varphi)^\dagger \tau^I \varphi \,.
\tag{2.3}
$$

The gauge field strength tensors are given by

$$
\begin{aligned}
G_{\mu\nu}^A &= \partial_\mu G_\nu^A - \partial_\nu G_\mu^A - g_s f^{ABC} G_\mu^B G_\nu^C \,, \\
W_{\mu\nu}^I &= \partial_\mu W_\nu^I - \partial_\nu W_\mu^I - g \epsilon^{IJK} W_\mu^J W_\nu^K \,, \\
B_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu \,,
\end{aligned}
\tag{2.4}
$$

while $\widetilde{X}_{\mu\nu} = \frac{1}{2} \epsilon_{\mu\nu\rho\sigma} X^{\rho\sigma}$, with $\epsilon_{0123} = +1$, denotes the dual tensor ($X = \{G^A, W^I, B\}$).

We consider a CP-conserving scalar potential [3]

$$V(\varphi_1, \varphi_2) = m_{11}^2\,\varphi_1^\dagger\varphi_1 + m_{22}^2\,\varphi_2^\dagger\varphi_2 - m_{12}^2\left(\varphi_1^\dagger\varphi_2 + \varphi_2^\dagger\varphi_1\right) + \frac{\lambda_1}{2}\left(\varphi_1^\dagger\varphi_1\right)^2 + \frac{\lambda_2}{2}\left(\varphi_2^\dagger\varphi_2\right)^2$$

$$+\lambda_3\,\varphi_1^\dagger\varphi_1\,\varphi_2^\dagger\varphi_2 + \lambda_4\,\varphi_1^\dagger\varphi_2\,\varphi_2^\dagger\varphi_1 + \frac{\lambda_5}{2}\left[\left(\varphi_1^\dagger\varphi_2\right)^2 + \left(\varphi_2^\dagger\varphi_1\right)^2\right] \tag{2.5}$$

with a Z_2 symmetry[3] (softly broken by a dimension-two term) preventing the existence of terms with odd powers of φ_1 and φ_2. Here all parameters are assumed to be real.[4] In order to study fluctuations around the VEVs that minimize the potential,[5] the two complex scalar fields are parameterized as

$$\varphi_a = \begin{pmatrix} \phi_a^+ \\ (v_a + \rho_a + i\eta_a)/\sqrt{2} \end{pmatrix}, \quad a = 1, 2\,. \tag{2.6}$$

The Lagrangian for the mass terms of the CP-odd (η_a), CP-even (ρ_a) and charged (ϕ_a^+) Higgses is

$$L_{M_H}^{(4)} = \frac{1}{2}\begin{pmatrix} \eta_1 \\ \eta_2 \end{pmatrix}^T m_\eta^2 \begin{pmatrix} \eta_1 \\ \eta_2 \end{pmatrix} + \begin{pmatrix} \phi_1^- \\ \phi_2^- \end{pmatrix}^T m_{\phi^\pm}^2 \begin{pmatrix} \phi_1^+ \\ \phi_2^+ \end{pmatrix} + \frac{1}{2}\begin{pmatrix} \rho_1 \\ \rho_2 \end{pmatrix}^T m_\rho^2 \begin{pmatrix} \rho_1 \\ \rho_2 \end{pmatrix} \tag{2.7}$$

with

$$m_\eta^2 = \left(v_1 v_2 \lambda_5 - m_{12}^2\right)\begin{pmatrix} -\dfrac{v_2}{v_1} & 1 \\ 1 & -\dfrac{v_1}{v_2} \end{pmatrix} \tag{2.8}$$

$$m_\rho^2 = \begin{pmatrix} \lambda_1 v_1^2 + m_{12}^2\dfrac{v_2}{v_1} & v_1 v_2\left(\lambda_3 + \lambda_4 + \lambda_5\right) - m_{12}^2 \\ v_1 v_2\left(\lambda_3 + \lambda_4 + \lambda_5\right) - m_{12}^2 & \lambda_2 v_2^2 + m_{12}^2\dfrac{v_1}{v_2} \end{pmatrix} \tag{2.9}$$

$$m_{\phi^\pm}^2 = \left[\dfrac{v_1 v_2}{2}\left(\lambda_4 + \lambda_5\right) - m_{12}^2\right]\begin{pmatrix} -\dfrac{v_2}{v_1} & 1 \\ 1 & -\dfrac{v_1}{v_2} \end{pmatrix} \tag{2.10}$$

where we eliminated m_{11}^2 and m_{22}^2 by the minimization conditions. The charged and CP-odd mass matrices have both one vanishing eigenvalue, which corresponds to the Goldstone bosons giving masses to the W and the Z, and a non-zero eigenvalue named $m_{H^\pm}^2$ and m_A^2, respectively. Both matrices are diagonalized by the same angle β defined as

$$\tan\beta \equiv v_2/v_1\,. \tag{2.11}$$

Another independent rotation angle, called α, enters the definition of the CP-even mass eigenstates h and H, with eigenvalues m_h and m_H respectively:

$$h = \rho_1 \sin\alpha - \rho_2 \cos\alpha\,,$$

$$H = -\rho_1 \cos\alpha - \rho_2 \sin\alpha \tag{2.12}$$

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[3]Note that assigning Peccei-Quinn charges to Higgs doublets and fermions has the same effect on the potential and the Yukawa couplings as imposing a Z_2 symmetry.

[4]For an analysis of the conditions for a CP-conserving Higgs sector we refer to [81].

[5]Since no CP violation is involved, both v_1 and v_2 can be taken to be real.

model	u_R	d_R	e_R
Type I	φ_2	φ_2	φ_2
Type II	φ_2	φ_1	φ_1
Lepton $-$ specific	φ_2	φ_2	φ_1
Flipped	φ_2	φ_1	φ_2

Table 2. Couplings of right-handed fermion singlets to Higgs doublets in 2HDMs with natural flavor conservation. These couplings can be enforced by an appropriate assignment of Z_2 (or Peccei-Quinn) charges to the scalar doublets and right-handed fermions.

where usually the lighter one is identified with the SM Higgs, with mass $m_h \approx 125\,\text{GeV}$.

Let us finally turn to the Yukawa part of the Lagrangian in eq. (2.1):

$$\mathcal{L}_Y = -Y_1^e\,\bar{l}\varphi_1 e - Y_2^e\,\bar{l}\varphi_2 e - Y_1^d\,\bar{q}\varphi_1 d - Y_2^d\,\bar{q}\varphi_2 d - Y_1^u\,\bar{q}\widetilde{\varphi}_1 u - Y_2^u\,\bar{q}\widetilde{\varphi}_2 u + h.c. \qquad (2.13)$$

Here we suppressed fermion flavor indices and defined $\widetilde{\varphi}^j = \epsilon_{jk}(\varphi^k)^\star$, using the totally antisymmetric ϵ_{jk} with $\epsilon_{12} = +1$. The Yukawa couplings $Y_{1,2}^f$ are understood to be 3×3 matrices in flavor space. If Y_1^f and Y_2^f are simultaneously non-zero, in general flavor changing neutral currents arise [82–84]. However, there are four 2HDMs with natural flavor conservation (see table 2) where only one of these couplings is present. As for CP-conservation in the potential, this can be achieved by an appropriate Z_2 charge assignment to right-handed fermions. When we discuss the extension to dimension six, we will assume that the terms are made Z_2-invariant in the same way.

3 Gauge invariant operators

In this section we list the independent gauge invariant operators up to dimension six in the 2HDM-EFT. They are defined before the EW symmetry breaking takes place, meaning that they are given in the interaction basis, as the mass basis is not yet defined. After EW symmetry breaking, the fermions acquire masses and also the Higgs mass matrices receive additional contributions compared to the 2HDM with dimension-four operators only.

At dimension five the generalization of the Weinberg operator reads

$$Q_{\nu\nu}^{11} = (\widetilde{\varphi}_1^\dagger l_p)^T C (\widetilde{\varphi}_1^\dagger l_r)\,, \quad Q_{\nu\nu}^{22} = (\widetilde{\varphi}_2^\dagger l_p)^T C (\widetilde{\varphi}_2^\dagger l_r) \qquad (3.1)$$

where C denotes the charge conjugation matrix.

The procedure we follow to obtain a complete set of independent operators at dimension six is the same as the one applied and thoroughly described in ref. [23] for SM-EFT. Obviously, the operators involving no Higgs doublets do not change compared to SM-EFT, and for them we refer the reader to ref. [23]. Using classical equations of motion, neglecting total derivatives, and imposing the constraint of vanishing total hypercharge, we derived a set of independent operators, which we classify like in the case of the SM-EFT as follows:

φ^6
$Q_\varphi^{111} = (\varphi_1^\dagger \varphi_1)^3$
$Q_\varphi^{112} = (\varphi_1^\dagger \varphi_1)^2(\varphi_2^\dagger \varphi_2)$
$Q_\varphi^{122} = (\varphi_1^\dagger \varphi_1)(\varphi_2^\dagger \varphi_2)^2$
$Q_\varphi^{222} = (\varphi_2^\dagger \varphi_2)^3$
$Q_\varphi^{(1221)1} = (\varphi_1^\dagger \varphi_2)(\varphi_2^\dagger \varphi_1)(\varphi_1^\dagger \varphi_1)$
$Q_\varphi^{(1221)2} = (\varphi_1^\dagger \varphi_2)(\varphi_2^\dagger \varphi_1)(\varphi_2^\dagger \varphi_2)$
$Q_\varphi^{(1212)1} = (\varphi_1^\dagger \varphi_2)^2(\varphi_1^\dagger \varphi_1) + h.c.$
$Q_\varphi^{(1212)2} = (\varphi_1^\dagger \varphi_2)^2(\varphi_2^\dagger \varphi_2) + h.c.$

Table 3. Operators in the 2HDM-EFT containing six Higgs doublets.

- φ^6: operators with Higgs doublets only (table 3), which modify the Higgs potential. We assumed that these operators respect the Z_2 symmetry present at dimension four.

- $\varphi^4 D^2$: operators with four Higgs doublets and two derivatives (table 4), which modify the kinetic terms of the Higgs fields, the Higgs-gauge boson interactions and the W and Z masses.

- $\Psi^2 \varphi X$: operators with two fermion fields, one field strength tensor and one Higgs doublet (table 5), which give rise to dipole interactions after EW symmetry breaking.

- $\varphi^2 X^2$: operators with two Higgs doublets and two field strength tensors (table 6).

- $\Psi^2 \varphi^2 D$: operators with two fermions fields, two Higgs doublets and one covariant derivative (table 7), which contribute to the fermion-Z and fermion-W couplings after EW symmetry breaking.

- $\Psi^2 \varphi^3$: operators containing two fermion fields and three Higgs doublets (table 8), which modify the relation between fermion masses and Higgs-fermion couplings.

As in the case of the dimension-four Lagrangian, we assume that the discrete Z_2 symmetry for operators involving Higgs and fermion fields is restored by an appropriate charge assignment to right-handed fermions. Concerning table 4, we stress that, as in the case of one Higgs doublet, operators with derivatives acting on two conjugated or two unconjugated fields are not independent.

4 Physical basis

In this section we discuss the modifications of the Higgs potential and the relation between the Yukawa couplings and the fermion masses induced by the dimension-six contributions.

$\varphi^4 D^2$		
$\Box$	φD	
$Q_\Box^{1(1)} = (\varphi_1^\dagger\varphi_1)\Box(\varphi_1^\dagger\varphi_1)$	$Q_{\varphi D}^{(1)11(1)} = \left[(D_\mu\varphi_1)^\dagger\varphi_1\right]\left[\varphi_1^\dagger(D^\mu\varphi_1)\right]$	$Q_{\varphi D}^{(1)21(2)} = \left[(D_\mu\varphi_1)^\dagger\varphi_2\right]\left[\varphi_1^\dagger(D^\mu\varphi_2)\right]+h.c.$
$Q_\Box^{2(2)} = (\varphi_2^\dagger\varphi_2)\Box(\varphi_2^\dagger\varphi_2)$	$Q_{\varphi D}^{(2)22(2)} = \left[(D_\mu\varphi_2)^\dagger\varphi_2\right]\left[\varphi_2^\dagger(D^\mu\varphi_2)\right]$	$Q_{\varphi D}^{(1)12(2)} = \left[(D_\mu\varphi_1)^\dagger\varphi_1\right]\left[\varphi_2^\dagger(D^\mu\varphi_2)\right]+h.c.$
$Q_\Box^{1(2)} = (\varphi_1^\dagger\varphi_1)\Box(\varphi_2^\dagger\varphi_2)$	$Q_{\varphi D}^{(1)22(1)} = \left[(D_\mu\varphi_1)^\dagger\varphi_2\right]\left[\varphi_2^\dagger(D^\mu\varphi_1)\right]$	$Q_{\varphi D}^{12(12)} = \left[\varphi_1^\dagger\varphi_2\right]\left[(D_\mu\varphi_1)^\dagger(D^\mu\varphi_2)\right]+h.c.$
$Q_\Box^{2(1)} = (\varphi_2^\dagger\varphi_2)\Box(\varphi_1^\dagger\varphi_1)$	$Q_{\varphi D}^{(2)11(2)} = \left[(D_\mu\varphi_2)^\dagger\varphi_1\right]\left[\varphi_1^\dagger(D^\mu\varphi_2)\right]$	$Q_{\varphi D}^{12(21)} = \left[\varphi_1^\dagger\varphi_2\right]\left[(D_\mu\varphi_2)^\dagger(D^\mu\varphi_1)\right]+h.c.$

Table 4. Operators with four Higgs doublets and two derivatives.

$\Psi^2\varphi X$		
G	W	B
$Q_{dG}^1 = (\bar{q}_p\sigma^{\mu\nu}T^A d_r)\varphi_1 G_{\mu\nu}^A$	$Q_{dW}^1 = (\bar{q}_p\sigma^{\mu\nu}d_r)\tau^I\varphi_1 W_{\mu\nu}^I$	$Q_{dB}^1 = (\bar{q}_p\sigma^{\mu\nu}d_r)\varphi_1 B_{\mu\nu}$
$Q_{dG}^2 = (\bar{q}_p\sigma^{\mu\nu}T^A d_r)\varphi_2 G_{\mu\nu}^A$	$Q_{dW}^2 = (\bar{q}_p\sigma^{\mu\nu}d_r)\tau^I\varphi_2 W_{\mu\nu}^I$	$Q_{dB}^2 = (\bar{q}_p\sigma^{\mu\nu}d_r)\varphi_2 B_{\mu\nu}$
$Q_{uG}^1 = (\bar{q}_p\sigma^{\mu\nu}T^A u_r)\widetilde{\varphi}_1 G_{\mu\nu}^A$	$Q_{uW}^1 = (\bar{q}_p\sigma^{\mu\nu}u_r)\tau^I\widetilde{\varphi}_1 W_{\mu\nu}^I$	$Q_{uB}^1 = (\bar{q}_p\sigma^{\mu\nu}u_r)\widetilde{\varphi}_1 B_{\mu\nu}$
$Q_{uG}^2 = (\bar{q}_p\sigma^{\mu\nu}T^A u_r)\widetilde{\varphi}_2 G_{\mu\nu}^A$	$Q_{uW}^2 = (\bar{q}_p\sigma^{\mu\nu}u_r)\tau^I\widetilde{\varphi}_2 W_{\mu\nu}^I$	$Q_{uB}^2 = (\bar{q}_p\sigma^{\mu\nu}u_r)\widetilde{\varphi}_2 B_{\mu\nu}$
	$Q_{eW}^1 = (\bar{l}_p\sigma^{\mu\nu}e_r)\tau^I\varphi_1 W_{\mu\nu}^I$	$Q_{eB}^1 = (\bar{l}_p\sigma^{\mu\nu}e_r)\varphi_1 B_{\mu\nu}$
	$Q_{eW}^2 = (\bar{l}_p\sigma^{\mu\nu}e_r)\tau^I\varphi_2 W_{\mu\nu}^I$	$Q_{eB}^2 = (\bar{l}_p\sigma^{\mu\nu}e_r)\varphi_2 B_{\mu\nu}$

Table 5. Operators containing two fermion fields, one Higgs doublet and a field strength tensor. Here $\sigma^{\mu\nu} = i\left[\gamma^\mu,\gamma^\nu\right]/2$.

$\varphi^2 X^2$	
GG, WW, BB	WB
$Q_{\varphi X}^{11} = (\varphi_1^\dagger\varphi_1)X_{\mu\nu}X^{\mu\nu}$	$Q_{\varphi WB}^{11} = (\varphi_1^\dagger\tau^I\varphi_1)W_{\mu\nu}^I B^{\mu\nu}$
$Q_{\varphi X}^{22} = (\varphi_2^\dagger\varphi_2)X_{\mu\nu}X^{\mu\nu}$	$Q_{\varphi WB}^{22} = (\varphi_2^\dagger\tau^I\varphi_2)W_{\mu\nu}^I B^{\mu\nu}$
$Q_{\varphi\widetilde{X}}^{11} = (\varphi_1^\dagger\varphi_1)\widetilde{X}_{\mu\nu}X^{\mu\nu}$	$Q_{\varphi\widetilde{W}B}^{11} = (\varphi_1^\dagger\tau^I\varphi_1)\widetilde{W}_{\mu\nu}^I B^{\mu\nu}$
$Q_{\varphi\widetilde{X}}^{22} = (\varphi_2^\dagger\varphi_2)\widetilde{X}_{\mu\nu}X^{\mu\nu}$	$Q_{\varphi\widetilde{W}B}^{22} = (\varphi_2^\dagger\tau^I\varphi_2)\widetilde{W}_{\mu\nu}^I B^{\mu\nu}$

Table 6. Operators with two scalar fields and two field strength tensors. X denotes G^A, W^I or B.

$\Psi^2 \varphi^2 D$	
(1)	(3)
$Q_{\varphi ud}^1 = i(\widetilde{\varphi}_1^\dagger i \overleftrightarrow{D}_\mu \varphi_1)(\overline{u}_p \gamma^\mu d_r)$	
$Q_{\varphi ud}^2 = i(\widetilde{\varphi}_2^\dagger i \overleftrightarrow{D}_\mu \varphi_2)(\overline{u}_p \gamma^\mu d_r)$	
$Q_{\varphi l}^{(1)1} = (\varphi_1^\dagger i \overleftrightarrow{D}_\mu \varphi_1)(\bar{l}_p \gamma^\mu l_r)$	$Q_{\varphi l}^{(3)1} = (\varphi_1^\dagger i \overleftrightarrow{D}_\mu^I \varphi_1)(\bar{l}_p \tau^I \gamma^\mu l_r)$
$Q_{\varphi l}^{(1)2} = (\varphi_2^\dagger i \overleftrightarrow{D}_\mu \varphi_2)(\bar{l}_p \gamma^\mu l_r)$	$Q_{\varphi l}^{(3)2} = (\varphi_2^\dagger i \overleftrightarrow{D}_\mu^I \varphi_2)(\bar{l}_p \tau^I \gamma^\mu l_r)$
$Q_{\varphi e}^1 = (\varphi_1^\dagger i \overleftrightarrow{D}_\mu \varphi_1)(\overline{e}_p \gamma^\mu e_r)$	
$Q_{\varphi e}^2 = (\varphi_2^\dagger i \overleftrightarrow{D}_\mu \varphi_2)(\overline{e}_p \gamma^\mu e_r)$	
$Q_{\varphi q}^{(1)1} = (\varphi_1^\dagger i \overleftrightarrow{D}_\mu \varphi_1)(\overline{q}_p \gamma^\mu q_r)$	$Q_{\varphi q}^{(3)1} = (\varphi_1^\dagger i \overleftrightarrow{D}_\mu^I \varphi_1)(\overline{q}_p \tau^I \gamma^\mu q_r)$
$Q_{\varphi q}^{(1)2} = (\varphi_2^\dagger i \overleftrightarrow{D}_\mu \varphi_2)(\overline{q}_p \gamma^\mu q_r)$	$Q_{\varphi q}^{(3)2} = (\varphi_2^\dagger i \overleftrightarrow{D}_\mu^I \varphi_2)(\overline{q}_p \tau^I \gamma^\mu q_r)$
$Q_{\varphi u}^1 = (\varphi_1^\dagger i \overleftrightarrow{D}_\mu \varphi_1)(\overline{u}_p \gamma^\mu u_r)$	
$Q_{\varphi u}^2 = (\varphi_2^\dagger i \overleftrightarrow{D}_\mu \varphi_2)(\overline{u}_p \gamma^\mu u_r)$	
$Q_{\varphi d}^1 = (\varphi_1^\dagger i \overleftrightarrow{D}_\mu \varphi_1)(\overline{d}_p \gamma^\mu d_r)$	
$Q_{\varphi d}^2 = (\varphi_2^\dagger i \overleftrightarrow{D}_\mu \varphi_2)(\overline{d}_p \gamma^\mu d_r)$	

Table 7. Operators in the 2HDM-EFT containing two fermions, two Higgs doublets and a covariant derivative. The superscripts (3), (1) label fermion bilinears transforming as an SU(2) triplet or singlet, respectively.

$\Psi^2 \varphi^3$		
e	d	u
$Q_{e\varphi}^{111} = (\bar{l}_p e_r \varphi_1)(\varphi_1^\dagger \varphi_1)$	$Q_{d\varphi}^{111} = (\overline{q}_p d_r \varphi_1)(\varphi_1^\dagger \varphi_1)$	$Q_{u\varphi}^{111} = (\overline{q}_p u_r \widetilde{\varphi}_1)(\varphi_1^\dagger \varphi_1)$
$Q_{e\varphi}^{122} = (\bar{l}_p e_r \varphi_1)(\varphi_2^\dagger \varphi_2)$	$Q_{d\varphi}^{122} = (\overline{q}_p d_r \varphi_1)(\varphi_2^\dagger \varphi_2)$	$Q_{u\varphi}^{122} = (\overline{q}_p u_r \widetilde{\varphi}_1)(\varphi_2^\dagger \varphi_2)$
$Q_{e\varphi}^{222} = (\bar{l}_p e_r \varphi_2)(\varphi_2^\dagger \varphi_2)$	$Q_{d\varphi}^{222} = (\overline{q}_p d_r \varphi_2)(\varphi_2^\dagger \varphi_2)$	$Q_{u\varphi}^{222} = (\overline{q}_p u_r \widetilde{\varphi}_2)(\varphi_2^\dagger \varphi_2)$
$Q_{e\varphi}^{211} = (\bar{l}_p e_r \varphi_2)(\varphi_1^\dagger \varphi_1)$	$Q_{d\varphi}^{211} = (\overline{q}_p d_r \varphi_2)(\varphi_1^\dagger \varphi_1)$	$Q_{u\varphi}^{211} = (\overline{q}_p u_r \widetilde{\varphi}_2)(\varphi_1^\dagger \varphi_1)$

Table 8. Operators in the 2HDM-EFT containing two fermion fields and three Higgs doublets.

The operators affecting Higgs kinetic terms and Higgs potential (and thus the Higgs mass matrices) are given in table 4 and table 3, respectively. The modified relations between fermion masses and Higgs-fermion couplings stem from the operators in table 8. The kinetic terms of the gauge boson fields receive contributions from the operators in table 6.

4.1 Kinetic terms, Higgs and gauge boson masses

The effect of the operators in table 4 on the kinetic terms of the Higgs fields amounts to

$$
\begin{aligned}
L_{H_{\mathrm{kin}}}^{(4)+(6)} = {} & \frac{1}{2}\begin{pmatrix}\partial_\mu\rho_1 \\ \partial_\mu\rho_2\end{pmatrix}^T
\begin{pmatrix}
1 + \dfrac{2\Delta_{\square}^{11}}{\Lambda^2} + \dfrac{\Delta_{\varphi D}^{11}}{2\Lambda^2} & \dfrac{\Delta_{\square}^{12}}{\Lambda^2} + \dfrac{\Delta_{\varphi D}^{12}}{2\Lambda^2} \\[2ex]
\dfrac{\Delta_{\square}^{12}}{\Lambda^2} + \dfrac{\Delta_{\varphi D}^{12}}{2\Lambda^2} & 1 + \dfrac{2\Delta_{\square}^{22}}{\Lambda^2} + \dfrac{\Delta_{\varphi D}^{22}}{2\Lambda^2}
\end{pmatrix}
\begin{pmatrix}\partial_\mu\rho_1 \\ \partial_\mu\rho_2\end{pmatrix} \\[2ex]
& + \frac{1}{2}\begin{pmatrix}\partial_\mu\eta_1 \\ \partial_\mu\eta_2\end{pmatrix}^T
\begin{pmatrix}
1 + \dfrac{\Delta_{\varphi D}^{11}}{2\Lambda^2} & \dfrac{\Delta_{\varphi D}^{12}}{2\Lambda^2} \\[2ex]
\dfrac{\Delta_{\varphi D}^{12}}{2\Lambda^2} & 1 + \dfrac{\Delta_{\varphi D}^{22}}{2\Lambda^2}
\end{pmatrix}
\begin{pmatrix}\partial_\mu\eta_1 \\ \partial_\mu\eta_2\end{pmatrix} \\[2ex]
& + \begin{pmatrix}\partial_\mu\phi_1^+ \\ \partial_\mu\phi_2^+\end{pmatrix}^\dagger
\begin{pmatrix}
1 & \dfrac{\Delta_{\varphi D}^+}{2\Lambda^2} \\[2ex]
\dfrac{\Delta_{\varphi D}^+}{2\Lambda^2} & 1
\end{pmatrix}
\begin{pmatrix}\partial_\mu\phi_1^+ \\ \partial_\mu\phi_2^+\end{pmatrix}
\end{aligned}
\tag{4.1}
$$

with

$$
\begin{aligned}
\Delta_{\square}^{11} &= -C_{\square}^{1(1)}v_1^2 \\
\Delta_{\square}^{12} &= -v_1 v_2\left(C_{\square}^{1(2)} + C_{\square}^{2(1)}\right) \\
\Delta_{\square}^{22} &= -C_{\square}^{2(2)}v_2^2 \\
\Delta_{\varphi D}^{11} &= C_{\varphi D}^{(1)11(1)}v_1^2 + C_{\varphi D}^{(1)22(1)}v_2^2 \\
\Delta_{\varphi D}^{22} &= C_{\varphi D}^{(2)11(2)}v_1^2 + C_{\varphi D}^{(2)22(2)}v_2^2 \\
\Delta_{\varphi D}^{12} &= v_1 v_2\left(C_{\varphi D}^{(1)12(2)} + C_{\varphi D}^{(1)21(2)} + C_{\varphi D}^{12(21)} + C_{\varphi D}^{12(12)}\right) \\
\Delta_{\varphi D}^+ &= v_1 v_2\left(C_{\varphi D}^{12(21)} + C_{\varphi D}^{12(12)}\right) .
\end{aligned}
\tag{4.2}
$$

The kinetic terms are made canonical by the shifts

$$
\begin{aligned}
\rho_1 &\to \rho_1\left(1 - \frac{\Delta_{\varphi D}^{11} + 4\Delta_{\square}^{11}}{4\Lambda^2}\right) - \left(\frac{\Delta_{\varphi D}^{12} + 4\Delta_{\square}^{12}}{4\Lambda^2}\right)\rho_2 \\[1ex]
\rho_2 &\to \rho_2\left(1 - \frac{\Delta_{\varphi D}^{22} + 4\Delta_{\square}^{22}}{4\Lambda^2}\right) - \frac{\Delta_{\varphi D}^{12} + 4\Delta_{\square}^{12}}{4\Lambda^2}\rho_1 \\[1ex]
\eta_1 &\to \eta_1\left(1 - \frac{\Delta_{\varphi D}^{11}}{4\Lambda^2}\right) - \frac{\Delta_{\varphi D}^{12}}{4\Lambda^2}\eta_2 \\[1ex]
\eta_2 &\to \eta_2\left(1 - \frac{\Delta_{\varphi D}^{22}}{4\Lambda^2}\right) - \frac{\Delta_{\varphi D}^{12}}{4\Lambda^2}\eta_1 \\[1ex]
\phi_1^+ &\to \phi_1^+ - \frac{\Delta_{\varphi D}^+}{4\Lambda^2}\phi_2^+ \\[1ex]
\phi_2^+ &\to \phi_2^+ - \frac{\Delta_{\varphi D}^+}{4\Lambda^2}\phi_1^+ .
\end{aligned}
\tag{4.3}
$$

The operators with gauge and Higgs fields in table 6 lead to the shifts

$$W_\mu^\pm \to W_\mu^\pm \left(1 + \frac{\Delta_{WW}}{\Lambda^2}\right) \tag{4.4}$$

$$Z_\mu \to Z_\mu \left(1 + \frac{\Delta_{ZZ}}{\Lambda^2}\right)$$

$$A_\mu \to A_\mu \left(1 + \frac{\Delta_{AA}}{\Lambda^2}\right) + Z_\mu \frac{\Delta_{AZ}}{\Lambda^2}$$

$$G_\mu^A \to G_\mu^A \left(1 + \frac{\Delta_{GG}}{\Lambda^2}\right)$$

with

$$\Delta_{GG} = v_1^2\, C_{\varphi G}^{11} + v_2^2\, C_{\varphi G}^{22}$$

$$\Delta_{WW} = v_1^2\, C_{\varphi W}^{11} + v_2^2\, C_{\varphi W}^{22}$$

$$\Delta_{ZZ} = c_w^2 \left(v_1^2 C_{\varphi W}^{11} + v_2^2 C_{\varphi W}^{22}\right) + s_w^2 \left(v_1^2 C_{\varphi B}^{11} + v_2^2 C_{\varphi B}^{22}\right) + c_w s_w \left(v_1^2 C_{\varphi WB}^{11} + v_2^2 C_{\varphi WB}^{22}\right)$$

$$\Delta_{AZ} = 2 c_w s_w \left[v_1^2 \left(C_{\varphi W}^{11} - C_{\varphi B}^{11}\right) + v_2^2 \left(C_{\varphi W}^{22} - C_{\varphi B}^{22}\right)\right] + \left(s_w^2 - c_w^2\right) \left(v_1^2 C_{\varphi WB}^{11} + v_2^2 C_{\varphi WB}^{22}\right)$$

$$\Delta_{AA} = s_w^2 \left(v_1^2 C_{\varphi W}^{11} + v_2^2 C_{\varphi W}^{22}\right) + c_w^2 \left(v_1^2 C_{\varphi B}^{11} + v_2^2 C_{\varphi B}^{22}\right) - c_w s_w \left(v_1^2 C_{\varphi WB}^{11} + v_2^2 C_{\varphi WB}^{22}\right)$$

where c_w and s_w denote cosine and sine of the weak mixing angle, respectively. These relations, together with the φD operators, affect the W and Z masses as

$$M_Z^2 = \frac{1}{4}\left(g^2 + g'^2\right)\left(v^2 + \frac{4v^2\Delta_{ZZ} + v_1^2\Delta_{\varphi D}^{11} + v_2^2\Delta_{\varphi D}^{22} + 2v_1 v_2 \Delta_{\varphi D}^{12}}{2\Lambda^2}\right) \tag{4.5}$$

$$M_W^2 = \frac{1}{4}g^2\left(v^2 + \frac{2v^2\Delta_{WW} + v_1 v_2 \Delta_{\varphi D}^+}{\Lambda^2}\right) \tag{4.6}$$

where $v^2 \equiv v_1^2 + v_2^2$.

The shift of the Higgs fields applied to the dimension-four potential and the φ^6 operators modify the Higgs mass terms in the Lagrangian. Choosing again to eliminate m_{11}^2 and m_{22}^2 via the minimization conditions, we obtained

$$L_{M_H}^{(4)+(6)} = \frac{1}{2}\begin{pmatrix}\eta_1 \\ \eta_2\end{pmatrix}^T \left(m_\eta^2 + \Delta m_\eta^2\right)\begin{pmatrix}\eta_1 \\ \eta_2\end{pmatrix} + \begin{pmatrix}\phi_1^- \\ \phi_2^-\end{pmatrix}^T \left(m_{\phi^\pm}^2 + \Delta m_{\phi^\pm}^2\right)\begin{pmatrix}\phi_1^+ \\ \phi_2^+\end{pmatrix}$$
$$+ \frac{1}{2}\begin{pmatrix}\rho_1 \\ \rho_2\end{pmatrix}^T \left(m_\rho^2 + \Delta m_\rho^2\right)\begin{pmatrix}\rho_1 \\ \rho_2\end{pmatrix} \tag{4.7}$$

with

$$\Delta m_\eta^2 = \Delta m_{\varphi D \eta}^2 + \Delta m_{\varphi^6 \eta}^2$$

$$\Delta m_\rho^2 = \Delta m_{\varphi D \rho}^2 + \Delta m_{\varphi^6 \rho}^2 \tag{4.8}$$

$$\Delta m_{\phi^\pm}^2 = \Delta m_{\varphi D \phi^\pm}^2 + \Delta m_{\varphi^6 \phi^\pm}^2$$

and

$$\Delta m_{\varphi D \phi^\pm}^2 = \left[(\lambda_4 + \lambda_5)\, v_1 v_2 - 2 m_{12}^2 \right] \frac{\Delta_{\varphi D}^+}{4\Lambda^2} \begin{pmatrix} -1 & \dfrac{v^2}{2 v_1 v_2} \\ \dfrac{v^2}{2 v_1 v_2} & -1 \end{pmatrix} \tag{4.9}$$

$$\left(\Delta m_{\varphi D \rho}^2 \right)_{11} = \frac{\left(v_1 \Delta_{\varphi D}^{12} - v_2 \Delta_{\varphi D}^{11} \right) m_{12}^2 - v_1^2 \left[v_1 \Delta_{\varphi D}^{11} \lambda_1 + v_2 \Delta_{\varphi D}^{12} (\lambda_3 + \lambda_4 + \lambda_5) \right]}{2 \Lambda^2 v_1}$$

$$\left(\Delta m_{\varphi D \rho}^2 \right)_{12} = \frac{\left[-v^2 \Delta_{\varphi D}^{12} + v_1 v_2 \left(\Delta_{\varphi D}^{22} + \Delta_{\varphi D}^{11} \right) \right] m_{12}^2}{4 \Lambda^2 v_1 v_2}$$

$$+ \frac{v_1 v_2 \left(-\Delta_{\varphi D}^{22} - \Delta_{\varphi D}^{11} \right) (\lambda_3 + \lambda_4 + \lambda_5) - \left(v_1^2 \lambda_1 + v_2^2 \lambda_2 \right) \Delta_{\varphi D}^{12}}{4 \Lambda^2}$$

$$\left(\Delta m_{\varphi D \rho}^2 \right)_{22} = \frac{\left(v_2 \Delta_{\varphi D}^{12} - v_1 \Delta_{\varphi D}^{22} \right) m_{12}^2 - v_2^2 \left[v_2 \Delta_{\varphi D}^{22} \lambda_2 + v_1 \Delta_{\varphi D}^{12} (\lambda_3 + \lambda_4 + \lambda_5) \right]}{2 \Lambda^2 v_2}$$

$$\Delta m_{\varphi D \eta}^2 = \frac{m_{12}^2 - v_1 v_2 \lambda_5}{2 \Lambda^2 v_1 v_2} \begin{pmatrix} v_1 v_2 \Delta_{\varphi D}^{12} - v_2^2 \Delta_{\varphi D}^{11} & \dfrac{v_1 v_2 \left(\Delta_{\varphi D}^{22} + \Delta_{\varphi D}^{11} \right) - v^2 \Delta_{\varphi D}^{12}}{2} \\ \dfrac{v_1 v_2 \left(\Delta_{\varphi D}^{22} + \Delta_{\varphi D}^{11} \right) - v^2 \Delta_{\varphi D}^{12}}{2} & v_1 v_2 \Delta_{\varphi D}^{12} - v_1^2 \Delta_{\varphi D}^{22} \end{pmatrix}$$

as well as

$$\Delta m_{\varphi^6 \eta}^2 = \Delta A \begin{pmatrix} -v_2/v_1 & 1 \\ 1 & -v_1/v_2 \end{pmatrix}$$

$$\Delta m_{\varphi^6 \phi^\pm}^2 = \Delta H^\pm \begin{pmatrix} -v_2/v_1 & 1 \\ 1 & -v_1/v_2 \end{pmatrix} \tag{4.10}$$

with

$$\Delta H^\pm = \frac{v_1 v_2}{2 \Lambda^2} \left[\left(C_\varphi^{(1212)1} + \frac{1}{2} C_\varphi^{(1221)1} \right) v_1^2 + \left(C_\varphi^{(1212)2} + \frac{1}{2} C_\varphi^{(1221)2} \right) v_2^2 \right]$$

$$\Delta A = \frac{v_1 v_2}{\Lambda^2} \left(C_\varphi^{(1212)1} v_1^2 + C_\varphi^{(1212)2} v_2^2 \right) \tag{4.11}$$

and

$$\left(\Delta m_{\varphi^6 \rho}^2 \right)_{11} = \frac{v_1^2}{\Lambda^2} \left(3 C_\varphi^{111} v_1^2 + \left(C_\varphi^{112} + 2 C_\varphi^{(1212)1} + C_\varphi^{(1221)1} \right) v_2^2 \right)$$

$$\left(\Delta m_{\varphi^6 \rho}^2 \right)_{12} = \frac{v_1 v_2}{\Lambda^2} \left[\left(C_\varphi^{112} + 2 C_\varphi^{(1212)1} + C_\varphi^{(1221)1} \right) v_1^2 + \left(C_\varphi^{122} + 2 C_\varphi^{(1212)2} + C_\varphi^{(1221)2} \right) v_2^2 \right]$$

$$\left(\Delta m_{\varphi^6 \rho}^2 \right)_{22} = \frac{v_2^2}{\Lambda^2} \left[\left(C_\varphi^{122} + 2 C_\varphi^{(1212)2} + C_\varphi^{(1221)2} \right) v_1^2 + 3 C_\varphi^{222} v_2^2 \right] . \tag{4.12}$$

Interestingly, the mass matrices for the CP-odd and the charged Higgs bosons are not diagonalized anymore by the angle β as in the case of the dimension-four potential. For this purpose, one needs two angles β_η and $\beta_{\phi^\pm}$, which do not satisfy $\tan \beta_{\eta, \phi^\pm} = v_1/v_2$.

This relation holds at dimension four but is broken by the φD operators. However, in the presence of $\Box$ and φ^6 operators only, the CP-odd and the charged Higgs mass matrices would still be diagonalized by β defined as $\tan\beta = v_1/v_2$. Also the angle α diagonalizing the CP-even mass matrix gets modified compared to the case of 2HDM without dimension-six operators. However, this effect can be accounted for by an appropriate redefinition of α, which is anyway a free parameter in the 2HDM.[6] The eigenvalues of the mass matrices also change, except those corresponding to the pseudo-Goldstone bosons, which are still zero at $\mathcal{O}(1/\Lambda^2)$, as we checked. Therefore, the effects on the eigenvalues can be absorbed into the definitions of m_h, m_H, $m_{H^\pm}$ and m_A.

4.2 Yukawa sector

After EW symmetry breaking, the fermion mass matrices in the presence of the dimension-six operators of table 8 are given by

$$m^f = \frac{v_1 Y_1^f}{\sqrt{2}} + \frac{v_2 Y_2^f}{\sqrt{2}} + \frac{1}{2\sqrt{2}\,\Lambda^2}\left(v_1^3\,C_{f\varphi}^{111} + v_1 v_2^2\,C_{f\varphi}^{122} + v_2^3\,C_{f\varphi}^{222} + v_1^2 v_2\,C_{f\varphi}^{211}\right), \qquad (4.13)$$

with $f = e, d, u$. Here m^f, $Y_{1,2}^f$ as well as the coefficients of the dimension-six operators are in general arbitrary 3×3 matrices in flavor space. However, since the eigenvalues of m^f are the physical fermion masses, by working in the basis where m^f is diagonal, the rotations that map onto this basis get implicitly absorbed in the definitions of $Y_{1,2}^f$ and $C_{f\varphi}$.

Depending on which version of the four 2HDMs with natural flavor conservation we are interested in (or which one is assumed to be the limiting case of the 2HDM of type III[7]), one can choose to eliminate either Y_1^f or Y_2^f from the Yukawa Lagrangian. Afterwards, it is straightforward to calculate the couplings of $\phi_{1,2}^+$, $\eta_{1,2}$ and $\rho_{1,2}$ to fermions.

Taking as an example the lepton Yukawa couplings and eliminating Y_1^e, we find

$$\begin{aligned}
\mathcal{L}_{Y^e}^{(6)} = {}& \frac{i}{v_1}\left(m^e - \frac{v_2}{2\sqrt{2}}\epsilon^e\right)\bar{l}P_R e\,\eta_1 + \frac{i}{2\sqrt{2}}\,\epsilon^e \bar{l}P_R e\,\eta_2 \\
& + \frac{1}{v_1}\left(\sqrt{2}\,m^e - \frac{v_2}{2}\epsilon^e\right)\bar{\nu}P_R e\,\phi_1^+ \\
& + \frac{1}{2}\,\epsilon^e \bar{\nu}P_R e\,\phi_2^+ + \left[\frac{C_{e\varphi}^{111}v_1^2 + C_{e\varphi}^{211}v_1 v_2}{\sqrt{2}\Lambda^2} + \frac{1}{v_1}\left(m^e - \frac{v_2\epsilon^e}{2\sqrt{2}}\right)\right]\bar{l}P_R e\,\rho_1 \\
& + \left(\frac{C_{e\varphi}^{222}v_2^2 + C_{e\varphi}^{122}v_1 v_2}{\sqrt{2}\Lambda^2} + \frac{\epsilon^e}{2\sqrt{2}}\right)\bar{l}P_R e\,\rho_2
\end{aligned} \qquad (4.14)$$

with

$$\epsilon^e = 2Y_2^e + \frac{C_{e\varphi}^{211}v_1^2 + C_{e\varphi}^{222}v_2^2}{\Lambda^2} \qquad (4.15)$$

---

[6]Nevertheless, one should keep in mind that in the limit of heavy Higgses A^0, H^0, $H^\pm$, α has to vanish as $\tan\beta \to \infty$ modulo corrections of order v^2/m_H^2. When dimension-six operators are included, this is still the case but additional (even smaller) corrections of order v^2/Λ^2 are present.

[7]Due to stringent flavor constraints (see ref. [85] for a recent analysis), a 2HDM with generic Yukawa couplings (i.e. of type III) should only slightly differ from one of the four 2HDMs with natural flavor conservation.

and P_R denoting the right-handed projector. The outcome in the case that Y_2^e is eliminated is obtained by simply interchanging $1 \leftrightarrow 2$. In order to express these couplings in terms of the physical Higgs fields, the rotations by (redefined) α, β_η and $\beta_{\phi^\pm}$ have to be applied. The analogous of eq. (4.14) for down-quarks has exactly the same structure. For up-type quarks all terms involving a charged or CP-odd Higgs switch sign and $\phi_{1,2}^+$ get replaced by $\phi_{1,2}^-$.

We can now point out an interesting effect arising in the 2HDM-EFT related to Higgs-pair production [86–89]. In the SM-EFT there is only one operator giving rise to fermion-fermion-Higgs-Higgs interactions and this is directly correlated to fermion-fermion-Higgs couplings [60, 90]. Therefore, its effect in the four-particle interaction is limited due to the constraints on the fermion-fermion-Higgs couplings. In our 2HDM-EFT three terms affect the fermion-fermion-Higgs couplings while the part of the Lagrangian relevant for Higgs pair production is

$$\frac{1}{2\sqrt{2}\Lambda^2}\overline{f}\Big(3C_{f\varphi}^{111}v_1\rho_1^2 + C_{f\varphi}^{122}v_1\rho_2^2 + C_{f\varphi}^{122}v_2\rho_1\rho_2 + C_{f\varphi}^{211}v_1\rho_1\rho_2 + C_{f\varphi}^{211}v_2\rho_1^2 + 3C_{f\varphi}^{222}v_2\rho_2^2\Big)f \,.$$
(4.16)

By comparing with eq. (4.14), it is clear that a cancellation between $C_{f\varphi}^{222}$ and $C_{f\varphi}^{122}$ could have quite different effects in fermion vertices with one or two Higgses. Therefore a suppressed modification to fermion-fermion-Higgs couplings accompanied by a sizable effect in pair production of SM-like Higgs bosons is possible in 2HDM-EFT.

5 Conclusions and outlook

In this article we extended the SM-EFT approach to NP to the case in which two Higgs doublets are dynamical degrees of freedom (2HDM-EFT). In this framework the effects of additional heavy particles (e.g. heavy SUSY partners in the MSSM) are parameterized in terms of higher-dimension operators and their Wilson coefficients. Our analysis enables systematic studies of the role played by additional degrees of freedom beyond the 2HDM field content. We derived a complete set of independent gauge-invariant dimension-six operators with two Higgs doublets of the same hypercharge under the assumption of a Z_2 symmetry involving the Higgs field and the right-handed fermions. These operators modify the Higgs potential and the relation between the fermion masses and Higgs-fermion-fermion couplings. We performed the transition to the physical basis by re-diagonalizing the Higgs, gauge boson and fermion kinetic terms and mass matrices. We showed that the CP-odd and charged Higgs mass matrices are in general not diagonalized by one angle β (defined by $\tan \beta = v_1/v_2$) but rather by two different angles. Finally, we derived the expressions for the Higgs-fermion-fermion couplings in the presence of dimension-six operators.

Even though a study of phenomenological applications of our framework is beyond the scope of this article, we pointed out one interesting example illustrating the differences between 2HDM-EFT and SM-EFT. In the latter, there is only one dimension-six operator giving rise to fermion-fermion-Higgs-Higgs interactions and its effect on pair production of SM Higgses is limited by constraints on Higgs-fermion-fermion couplings. In the 2HDM-EFT instead, there are three operators entering Higgs-fermion-fermion couplings, whose contributions can cancel each other. The same operators, as well as additional ones, also con-

tribute to the couplings of Higgs pairs to fermions but with different prefactors. Higgs pair production is decoupled from single-Higgs-fermion interactions and could thus be sizable.

Acknowledgments

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