Global constraints on Yukawa operators in the standard model effective theory

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Abstract: CP-violating contributions to Higgs-fermion couplings are absent in the standard model of particle physics (SM), but are motivated by models of electroweak baryogenesis. Here, we employ the framework of the SM effective theory (SMEFT) to parameterise deviations from SM Yukawa couplings. We present the leading contributions of the relevant operators to the fermionic electric dipole moments (EDMs). We obtain constraints on the SMEFT Wilson coefficients from the combination of LHC data and experimental bounds on the electron, neutron, and mercury EDMs. We perform, for the first time, a combined fit to LHC and EDM data allowing the presence of CP-violating contributions from several fermion species simultaneously. Among other results, we find non-trivial correlations between EDM and LHC constraints even in the multi-parameter scans, for instance, when floating the CP-even and CP-odd couplings to all third-generation fermions.

Keywords: SMEFT, CP Violation, Electric Dipole Moments

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1 Introduction

Charge-Parity (CP) violating contributions to Higgs-fermion couplings are a well-motivated possibility of physics beyond the standard model (SM) that might help address the problem of baryogenesis with new dynamics at the electroweak scale. As is well-known, any such contributions are strongly constrained by null measurements of electric dipole moments (EDMs), and less strongly by Higgs production and decay data from colliders. It is also well-known that the presence of several CP-violating phases can lead to cancellations and thus weaker bounds [1]. The interplay of EDM and collider constraints, in the presence of several CP-violating Higgs couplings to fermions simultaneously, is less well-known. In ref. [2] constraints in the presence of two CP-violating parameters have been studied in detail. In this work, we scan over up to six different parameters.
Frequently, model-independent bounds on such interactions have been obtained in the so-called “κ framework” by rescaling the SM Yukawa coupling by an overall factor and a complex phase. See, e.g., refs. [3–5] for recent and future analyses of LHC data. In ref. [6], bounds on these factors have been obtained by studying contributions to EDMs through Barr-Zee diagrams with internal fermion loops. In refs. [7, 8] it was shown by explicit calculation within the κ framework that bosonic two-loop diagrams have a large impact on the EDM bounds for light quarks. However, it was later pointed out that the way these bosonic contributions had been computed lead to a gauge dependent result [9]. This is related to the fact that the naive implementation of the κ framework, in which only the dimension-four Higgs couplings are modified, is not a consistent quantum field theory. A consistent alternative that resembles the κ framework most closely would be to use Higgs effective theory (HEFT) [10]. In HEFT, the dimension-four couplings are complemented by the necessary higher dimension interactions to facilitate gauge-independent results.

Another consistent way of obtaining model-independent bounds for CP-violating Higgs-fermion interactions is the use of the SM effective field theory (SMEFT) [11]; see refs. [2, 12–14] for recent work in this direction. We take the same approach in this article. For the purpose of this work, we assume that the contributions to the EDMs of the SM fermions arise from physics at a scale Λ that is sufficiently higher than the electroweak scale. This permits us to parameterise deviations from the SM in terms of operators in SMEFT. We will assume that any of the operators of the form \((H^T H) \bar{F}_L f_R H\) may have non-zero coefficients. Here, \(H\) is the complex Higgs doublet field, \(F_L\) a left-handed doublet fermion field, \(f_R\) a right-handed singlet fermion field, and we have suppressed flavour indices. The reason for focusing on this class of operators is that they are the only dimension-six operators that induce tree-level modifications to Higgs-fermion couplings. See section 2 for a detailed discussion. These operators will contribute to leptonic and hadronic dipole moments via a series of matching and renormalization-group (RG) evolution. Our analysis takes into account the leading effects that arise at two-loop order in the electroweak interactions, as well as leading-logarithmic QCD corrections below the electroweak scale. We neglect any effects of CP-odd, flavour off-diagonal operators. The study of these effects is relegated to future work.

To understand the complex interplay between the operators we consider requires us to vary multiple Wilson coefficients simultaneously. To explore this computationally challenging multidimensional parameter space, we make use of the GAMBIT [15] global fitting framework, to which we have added a new module that allows for calculations of EDMs in the SMEFT, as well as the corresponding experimental likelihoods. We have also expanded the existing ColliderBit [16] module to be able to determine constraints on the Wilson coefficients from measured properties of the Higgs boson at the LHC.

This work contains several new aspects and presents nontrivial results. For the first time, we perform a multi-parameter fit to both LHC and EDM data. This corresponds to a more realistic scenario than single-parameter fits, as in a UV-complete model several CP phases are expected to be present. Moreover, we complement and correct some of the analytic expressions in the literature. In addition to updating existing constraints on either individual CP-odd Yukawa couplings (figures 4 and 5), or two CP-odd Yukawa
couplings as in ref. [2] (figure 6), we scan over up to six CP-even/CP-odd Yukawa couplings simultaneously for the first time (figures 7–11). While this represents still only a subset of all Yukawa operators, we argue that with current experimental results, including more parameters will not lead to additional significant constraints. For instance, allowing for any CP-odd contribution to the Yukawa of a light fermion (electron, or up and down quark) would allow to fully cancel the corresponding EDM constraint, thus leaving only the LHC bounds in a trivial way. However, we find several nontrivial effects. Focusing on the heavier fermions, we find an intricate interplay between LHC and precision constraints that weakens the EDM bounds without lifting them fully. Not even by allowing for six independent parameters can we cancel all EDM constraints among the third-generation fermions, and nontrivial EDM bounds remain. A detailed discussion of these issues is found in section 7.

In our analysis, we do not take into account perturbative and hadronic uncertainties. The effects of hadronic uncertainties on EDM bounds have been studied in detail in ref. [2] (see also ref. [17]), and will not be repeated here. They are relevant mainly for bounds on CP-violating Higgs couplings to the light quarks that arise from hadronic EDMs, as collider bounds are nearly absent. When allowing for the presence of several Wilson coefficients at the same time, EDM bounds tend to get cancelled and only collider bounds remain, so hadronic uncertainties are less relevant. Note also that bounds from arising from the electron EDM have no hadronic uncertainty. However, there is a large, previously overlooked perturbative uncertainty regarding the bottom and charm couplings, as discussed in section 3. This uncertainty is analogous to the case discussed in ref. [18], and will be studied in detail in a forthcoming publication [19]. As not all relevant higher-order corrections are currently known, they are neglected in this analysis.

This paper is organised as follows. In section 2 we specify the operators that we want to constrain; namely, those that modify the SM Yukawa couplings at tree level. We then derive the modified Higgs-fermion couplings in the broken electroweak phase. In section 3 we discuss the effective theory valid below the electroweak scale. The analytic contributions of the SMEFT operators to the partonic EDM of leptons and quarks are collected in section 4. Most of these results are taken from the literature, although a few results are presented here for the first time. Furthermore, we correct several errors in the literature. In section 5 we summarize the contributions of the partonic EDMs to the EDM of the physical systems that are actually constrained by experiment. Section 6 describes the collider constraints on the SMEFT operators, and contains a short discussion of the expected contribution of operator with mass dimension eight. Our main results are contained in section 7. We conclude in section 8. Appendix A contains the full generic $R_\xi$-gauge Lagrangian in the broken phase. In appendix B we discuss the alternative flavour basis in the electroweak broken phase that has been used in ref. [20].

2 Effective theory above the electroweak scale — SMEFT

In this work we consider the SM augmented with SMEFT operators that induce tree-level modifications to the Yukawa couplings. In the unbroken phase, the relevant part of the
Lagrangian reads

\[ \mathcal{L}_{\text{Yukawa}} = -\bar{Q}_L \tilde{H} Y_u u_R + \frac{1}{\Lambda^2} (H^\dagger H) \bar{Q}_L \tilde{H} C'_{uH} u_R \\
- \bar{Q}_L \tilde{H} Y_d d_R + \frac{1}{\Lambda^2} (H^\dagger H) \bar{Q}_L \tilde{H} C'_{dH} d_R \\
- \bar{L}_L \tilde{H} Y_\ell \ell_R + \frac{1}{\Lambda^2} (H^\dagger H) \bar{L}_L \tilde{H} C'_{\ell H} \ell_R + \text{h.c.} \quad (2.1) \]

Here, $u_R$, $d_R$, $\ell_R$ denote the triplets (in generation space) of right-handed up-, down-, and charged-lepton fields, while $Q_L$ and $L_L$ are the corresponding left-handed triplets. In accordance, the Yukawa matrices $Y_u$, $Y_d$, $Y_\ell$ and Wilson coefficients $C'_{uH}$, $C'_{dH}$, $C'_{\ell H}$, are generic complex $3 \times 3$ matrices. The primes indicate that they are not necessarily couplings in the mass-eigenstate basis. Phases in the Yukawa couplings and Wilson coefficients can potentially induce beyond-the-SM (BSM) CP-violating contributions to Higgs-fermion couplings. As is well-known, not all these complex parameters are physical, due to the freedom of choosing the phases of the left- and right-handed fermion fields. To isolate the physical BSM parameters, we express as many parameters as possible in terms of observed quantities, e.g., the fermion masses and CKM matrix elements. To this end, we rewrite the Lagrangian in the broken phase using the linear decomposition of the Higgs field as

\[ H = \left( \frac{1}{\sqrt{2}} (v + h + i G^0) \right) \quad (2.2) \]

The terms that induce the fermion masses are

\[ \mathcal{L}_{\text{mass}} = -\sum_{f=u,d,\ell} \frac{v}{\sqrt{2}} \bar{f}_L \left( Y_f - \frac{v^2}{2\Lambda^2} C'_{fH} \right) f_R + \text{h.c.} \quad (2.3) \]

In analogy to the SM, we diagonalise the complete matrix in parentheses by biunitary transformations, parameterised by the field redefinitions

\[ f_L \rightarrow U_f f_L, \quad f_R \rightarrow W_f f_R, \quad (2.4) \]

with $f = u, d, \ell$, and $U_f$, $W_f$ complex $3 \times 3$ matrices chosen such that after this rotation, the mass Lagrangian is

\[ \mathcal{L}_{\text{mass}} = -\sum_{f=u,d,\ell} \frac{v}{\sqrt{2}} \bar{f}_L y_f^{\text{SM}} f_R + \text{h.c.} \quad (2.5) \]

Here, the $y_f^{\text{SM}}$ are real and diagonal matrices with entries that correspond to the observed fermion masses, i.e., $m_f = \frac{v}{\sqrt{2}} y_f^{\text{SM}}$. In other words, the conditions on $U_f$, $W_f$ read

\[ y_f^{\text{SM}} \equiv U_f^\dagger \left( Y_f - \frac{v^2}{2\Lambda^2} C'_{fH} \right) W_f. \quad (2.6) \]

The kinetic terms of the fermions are also affected by the transformation in eq. (2.4), which leads to the appearance of the CKM matrix $V_{\text{CKM}} \equiv U_u^\dagger W_d$ in the charged gauge interactions of the left-handed quarks, in complete analogy to the SM.
The interaction terms of the fermions with a single Higgs field from eq. (2.1) after performing the transformation in eq. (2.4) are given by

$$\mathcal{L}_h = - \sum_{f=u,d,\ell} \frac{h}{\sqrt{2}} \bar{f}_L U^\dagger_f \left( Y_f - \frac{3v^2}{2\Lambda^2} C'_{fH} \right) W_f f_R + \text{h.c.}$$

where we defined $C_{fH} \equiv U^\dagger_f C'_{fH} U_f$. We see that, in the most general case, we can parametrise any deviation with respect to the SM by the rotated Wilson coefficients $C_{fH}$. For later convenience, we split these coefficients into their real and imaginary parts,

$$C_{fH,ij} = \text{Re}[C_{fH,ij}] + i \text{Im}[C_{fH,ij}] \equiv C_{fH,+ij} + iC_{fH,-ij}.$$  (2.8)

with $C_{fH,\pm ij}$ real. This implies that the operator combinations proportional to the diagonal terms $C_{fH,\pm ii}$ are self-adjoint.

In the current work we consider the effect of the flavour-diagonal contributions, $C_{fH,ii}$, for the case that they contain additional sources of CP violation. This implies, as we shall see, that $C_{fH,-ii}$ is non-zero. Note, however, that in the absence of any type of flavour alignment with the SM Yukawas these operators also induce flavour violation beyond the SM, which we do not consider in this work. In a concrete model, we thus expect additional to constraints from such flavour off-diagonal contributions of $C_{fH}$ (see refs. [21, 22] for a recent discussion). It is, however, possible to obtain both CP-even and CP-odd modifications in the flavour-diagonal Yukawas without any flavour-violation beyond the SM. This corresponds to a UV scenario in which the unitary transformations in eqs. (2.6) and (2.4) that rotate to the mass eigenstates simultaneously diagonalise the coefficients $C'_{fH}$. In this setup $C_{fH}$ is diagonal but not necessarily real, meaning that CP violation beyond the SM is still possible.

In unitarity gauge, all BSM effects in our setup are contained in the Yukawa Lagrangian in the mass-eigenstate basis

$$\mathcal{L}_{\text{Yukawa}} = \sum_f \left[ \left( -m_f - m_f \frac{h}{v} + \frac{v^3}{2\sqrt{2}\Lambda^2} C_{fH,+} \left( 2\frac{h}{v} + 3\frac{h^2}{v^2} + \frac{h^3}{v^3} \right) \right) \bar{f} f + \frac{v^3}{2\sqrt{2}\Lambda^2} C_{fH,-} \left( 2\frac{h}{v} + 3\frac{h^2}{v^2} + \frac{h^3}{v^3} \right) \bar{f} i\gamma_5 f \right].$$  (2.9)

Here, the sum runs over all charged fermion fields, $f = u, d, s, c, b, t, e, \mu, \tau$, and we have traded the generation indices on the Wilson coefficients in favour of the flavour label, $f$, since we focus on the flavour-diagonal case. We will use this more compact notation in the remainder of this work. In appendix A we present the corresponding Lagrangian for generic $R_\xi$ gauge. A different basis has been used in ref. [20] to present the bounds, see appendix B.
Often, the $\kappa$ framework is employed to parametrise flavour-diagonal CP-violation in the Yukawas [6, 8, 18]. The corresponding Lagrangian reads
\[
\mathcal{L}_{\text{Yukawa},\kappa} = -\sum_f \frac{y_f^{\text{SM}}}{\sqrt{2}} \kappa_f h \bar{f} \left( \cos \phi_f + i \gamma_5 \sin \phi_f \right) f,
\]
(2.10)
with $\kappa_f$ a real parameter controlling the absolute value of the modification and $\phi_f \in [0, 2\pi)$ a CP-violating phase, such that $\kappa_f^{\text{SM}} = 1$ and $\phi_f^{\text{SM}} = 0$ reproduces the SM. Eq. (2.10) should be thought of as the dimension-four part of the corresponding HEFT Lagrangian in unitarity gauge. As long dimension-six SMEFT is a good approximation, and the $h^2$ and $h^3$ interactions in eq. (2.9) do not affect the observables considered, the bounds on $C_{fH}^\pm$ directly translate to bounds on the $\kappa$-framework parameters via
\[
\kappa_f \cos \phi_f \approx 1 - \frac{v}{\sqrt{2} m_f} \Lambda^2 C_{fH^+}, \quad \kappa_f \sin \phi_f \approx -\frac{v}{\sqrt{2} m_f} \Lambda^2 C_{fH^-}.
\]
(2.11)
This is actually the case for the observables and the precision that we consider. A possible exception is the top contribution to the LHC observables, see the discussion in section 6.

Using the Lagrangian in eq. (2.9) we will be able to set constraints on the $C_{fH}^\pm$ couplings from LHC measurements. Low-energy probes of CP-violation, however, also place significant constraints on these couplings. Our aim is to capture the numerically leading effects. To this end we work with a tower of effective theories. Apart from a few exceptions that we discuss below, the RG evolution from $\Lambda$ to $\mu_{\text{ew}}$ within SMFET can be neglected as (most) Yukawa operators do not mix into the CP-violating dipoles for the electron and light-quarks, relevant for EDMs. In this case the leading effects are obtained at the electroweak scale by matching to the effective theory in which the heavy degrees of the SM are integrated out, and by subsequently running to the hadronic scale. In contrast, the few contributions that are induced from (two-loop) mixing within SMEFT will come with a UV-sensitive logarithm, $\log M_h/\Lambda$.

3 Effective theory below the electroweak scale

We will extract constraints on the coefficients in eq. (2.9) from the experimental bounds on the electron, neutron, and mercury EDMs. We will make the assumption that these EDMs are induced solely by the electron and partonic CP-violating electric and chromoelectric dipole operators, and the purely gluonic CP-violating Weinberg operator [23] with coefficients $d_e, d_q, \tilde{d}_q, w$, respectively. Therefore, we neglect (subleading) contributions from the matrix elements of four-fermion operators. Here, $q$ collectively denotes light quarks, $u, d, s$. The charm-quark is typically not included since presently there are no lattice computations for the hadronic matrix elements of its operators. The above coefficients are traditionally defined via the effective, CP-odd Lagrangian valid at hadronic energies $\mu_{\text{had}} \simeq 2$ GeV [24],
\[
\mathcal{L}_{\text{eff}} = -d_e \frac{i}{2} \bar{e} \sigma^{\mu\nu} \gamma_5 e F^\mu_{\nu} - d_q \frac{i}{2} \bar{q} \sigma^{\mu\nu} q F^\mu_{\nu} - \tilde{d}_q \frac{i g_s}{2} \bar{q} \sigma^{\mu\nu} T^a q G^a_{\mu\nu} + \frac{1}{3} w f_{abc} G^a_{\mu\nu} G^b_{\nu\sigma} \tilde{G}^{c,\mu\nu}.
\]
(3.1)
Here, $\sigma^{\mu\nu} = \frac{i}{2} [\gamma^\mu, \gamma^\nu]$, $T^a$ are the generators of SU($n_c$) in the fundamental representation normalised as $\text{Tr} [T^a, T^b] = \delta^{ab}/2$, $n_c = 3$ is the number of colours, and we collect our conventions for the field strength and its dual below. The contributions from the Weinberg operator turn out to be subdominant because of its small nuclear matrix elements [24, 25], but we keep them for completeness.

The partonic dipole moments $d_q$ and $\tilde{d}_q$ as well as the coefficient $w$ are obtained from the SMEFT Wilson coefficients in eq. (2.9) via a sequence of matching at the weak scale, and the bottom- and charm-flavour thresholds, as well as the RG evolution between each scale, as outlined in refs. [6, 18, 26–28]. The resummation of logarithms is phenomenologically relevant for the hadronic dipole operators that mix under QCD. In the current analysis, we work at the leading-logarithmic order for quark operators. QCD corrections are also large for the bottom- and charm-quark contributions to the electron EDM. They are currently unknown [19] and are therefore not included in our analysis.

To perform the sequence of matching and RG evolution connecting the electroweak-scale Lagrangian, eq. (2.9), with the hadronic-scale one, eq. (3.1), requires the full flavour-conserving, CP-odd effective Lagrangian below the electroweak scale up to mass dimension six. In the conventions and basis of ref. [18], adapted from ref. [27], it reads:

$$
\mathcal{L}_{\text{eff}} = -\sqrt{2} G_F \left( \sum_{q \neq q'} C_i^{qq'} O_i^{qq'} + \frac{1}{2} \sum_{i=3,4} C_i^{qq'} O_i^{qq'} \right) + \sum_q \sum_{i=1,\ldots,4} C_i^q O_i^q + C_w O_w + \sum_\ell C_3^\ell O_3^\ell ,
$$

where the sums run over all active quarks with masses below the weak scale, e.g., $q, q' = u, d, s, c, b$ in the five-flavour theory. The linearly independent operators are\(^1\)

$$
O_1^{qq'} = (\bar{q} q) (\bar{q}' i\gamma_5 q') , \quad O_2^{qq'} = (\bar{q} T^a q) (\bar{q}' i\gamma_5 T^a q') ,
$$

$$
O_3^{qq'} = \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} (\bar{q} \sigma_{\mu\rho} q) (\bar{q}' \sigma_{\nu\sigma} q') , \quad O_4^{qq'} = \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} (\bar{q} \sigma_{\mu\rho} T^a q) (\bar{q}' \sigma_{\nu\sigma} T^a q') ,
$$

$$
O_5^{q} = (\bar{q} q) (\bar{q} i\gamma_5 q) , \quad O_6^{q} = \frac{1}{2} \epsilon^{\mu\nu} (\bar{q} \sigma_{\mu\nu} T^a q) (\bar{q} \sigma_{\mu\nu} T^a q) ,
$$

$$
O_7^{q} = \frac{e Q_q m_q}{2 g_s^2} \bar{q} \sigma_{\mu\nu} q \tilde{F}_{\mu\nu} , \quad O_8^{q} = -\frac{1}{2} m_q \bar{q} \sigma_{\mu\nu} T^a q \tilde{C}_\mu^a ,
$$

$$
O_w = -\frac{1}{3} \frac{m_q}{g_s} \epsilon^{abc} G_{\mu\sigma}^a C_{\nu\rho}^b \tilde{G}^{c\mu\nu} .
$$

\(^1\)The definition of the $O_i^{qq'}$, $O_i^q$ operators in terms of the $\epsilon$-tensor is convenient when going beyond leading-logarithmic approximation, as they remain selfadjoint also in $d = 4 - 2\epsilon$ dimensions. It is thus convenient to also define $O_2^q$ and the dipoles $O_7^q$ and $O_8^q$ in an analogous manner. To express the operators in the more conventional $\sigma^{\mu\nu} a^5$ notation one can use the $d = 4$ relation

$$
\frac{d}{4} \epsilon^{\mu\nu\rho\sigma} = \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} \sigma_{\rho\sigma} ,
$$

where $\epsilon_{0123} = -\epsilon^{0123} = 1$. For details concerning the next-to-leading-logarithmic analysis see ref. [18].
The powers of $e$ and $g_s$ in the operator definitions above are fixed such that all Wilson coefficients have the same $h$ dimension, thus making the calculation of operator mixing more transparent.

$C_w$ as well as some of the anomalous dimensions which control operator mixing depend on the conventions for the covariant derivatives for quarks and leptons, the field strength tensor, and the dual tensors. Our conventions (same as in ref. [18]) read

$$D_\mu^a \equiv \partial_\mu - ig_s T^a G^a_{\mu} + ieQ_q A_\mu, \quad D_\mu^f \equiv \partial_\mu + ieQ_f A_\mu,$$

$$C_{\mu\nu}^a \equiv \partial_\mu C_{\nu}^a - \partial_\nu C_{\mu}^a + g_s f^{abc} G^b_{\mu} G^c_{\nu}, \quad \tilde{G}^{a,\mu\nu} \equiv \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} G^a_{\rho\sigma}, \quad \tilde{F}^{\mu\nu} \equiv \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} F_{\rho\sigma},$$

with $\epsilon_{0123} = -\epsilon^{0123} = 1$.

The comparison of eqs. (3.1) and (3.2) gives the relation between $d_e, d_q, \tilde{d}_q, w$ and the Wilson coefficients of the three-flavour EFT evaluated at the hadronic scale $\mu_{\text{had}} = 2 \text{ GeV}$:

$$d_e = \sqrt{2} G_F \frac{e Q_e}{4\pi\alpha} m_e C_3^e,$$

$$d_q(\mu_{\text{had}}) = \sqrt{2} G_F \frac{e Q_q}{4\pi\alpha_s} m_q C_3^q(\mu_{\text{had}}),$$

$$\tilde{d}_q(\mu_{\text{had}}) = -\sqrt{2} G_F \frac{1}{4\pi\alpha_s} m_q C_4^q(\mu_{\text{had}});$$

$$w(\mu_{\text{had}}) = \sqrt{2} G_F \frac{1}{g_s} C_w(\mu_{\text{had}}).$$

Note that $C_3^e$ is scale independent with regard to the strong interaction, i.e., $C_3^e(\mu_{\text{had}}) = C_3^e(\mu_{\text{ew}})$ with $\mu_{\text{ew}} \approx 100 \text{ GeV}$. The procedure to calculate $C_3^e$, $C_3^q(\mu_{\text{had}})$, $C_4^q(\mu_{\text{had}})$, and $C_w(\mu_{\text{had}})$ depends on the flavour quantum numbers of the SMEFT operators. In the following section, we describe the different cases and present the corresponding results.

### 4 Electroweak matching and RG evolution to the hadronic scale

In this section we discuss how to obtain the Wilson coefficients of the EFT below the electroweak scale as a function of the Yukawa SMEFT Wilson coefficients. Since all operators in eq. (3.2) are CP odd, all contributions are proportional to the CP-odd couplings $C_{fH-}$. The final results for the required initial conditions $C_{1fH-}^{\mu_{\text{ew}}}$, $C_{2fH-}^{\mu_{\text{ew}}}$, $C_{wH-}^{\mu_{\text{ew}}}$, $C_{3fH-}^{\mu_{\text{ew}}}$, and $C_{4fH-}^{\mu_{\text{ew}}}$ are collected below in eqs. (4.3), (4.4), (4.5), (4.6), and (4.7), respectively. These are all the initial conditions required for the leading-log QCD analysis.

We begin the discussion with contributions that are UV sensitive, i.e., proportional to $\log(\mu_{\text{ew}}/\Lambda)$. There are two equivalent ways of obtaining them. One way is to identify those SMEFT Yukawa operators $O_{fH}$ that mix into the SMEFT dipole operators $O_{fW}$, $O_{fB}$ and $O_{qG}$. The mixing induces dipole Wilson coefficient proportional to the leading UV logarithm $\log(\mu_{\text{ew}}/\Lambda)$, and the subsequent tree-level matching onto the Lagrangian eq. (3.2) potentially induces non-zero coefficients $C_{3fH-}^{\mu_{\text{ew}}}$, $C_{3fH-}^{\mu_{\text{ew}}}$, and $C_{4fH-}^{\mu_{\text{ew}}}$. In fact, only the following cases occur: the operator $O_{eH}$ mixes at the two-loop, electroweak level into the leptonic dipole operators

$$O_{eW} = \tilde{L}_L^\sigma \bar{\sigma} \epsilon_R H W_{\mu\nu}^a, \quad O_{eB} = \tilde{L}_L^\sigma \epsilon_R H B_{\mu\nu}. \quad (4.1)$$
Analogously, the Yukawa operators $O_{qH}$ with $q = u, d, s$ mix into the quark dipole operators

$$O_{qW} = \bar{Q}_L \sigma^a \sigma^{\mu\nu} q_R \hat{H} W_{\mu\nu}^a, \quad O_{qB} = \bar{Q}_L \sigma^{\mu\nu} q_R \hat{H} B_{\mu\nu},$$

with $\hat{H} \equiv H$ for $q = d, s$ and $\hat{H} \equiv \tilde{H}$ for $q = u$. Here, $\sigma^a$ are the Pauli matrices acting in SU(2)$_L$ space. (Here, we considered only mixing into dipole operators with light fermions.) The two-loop mixing has been calculated for the electron case in ref. [28].

An alternative way of obtaining these $\log \mu_{\text{ew}} / \Lambda$ terms is to directly perform the two-loop electroweak matching calculation to extract $C_{fH}^3(\mu_{\text{ew}})$, for the light-fermion flavours $f = e, u, d, s$ (see representative diagrams in the lower panels in figure 2). For these coefficients, the contributions to the matching proportional to the SMEFT coefficients $C_{fH-}$ are divergent and thus include a term proportional to the UV-sensitive logarithm $\log M_H / \Lambda$. We performed the matching calculation explicitly and find that only the contributions to $C_{fH}^3(\mu_{\text{ew}})$ are UV sensitive, while the matching for $C_{qH}^4(\mu_{\text{ew}})$ is UV finite. To obtain gauge-independent results in this calculation, it was necessary [9] to include the dimension-five vertices in the Lagrangian in the broken phase (see appendix A) that were missed in ref. [7]. The UV-divergent contributions are a direct consequence of including these dimension-five vertices. The corresponding logarithmic terms in our results for $C_{fH}^3(\mu_{\text{ew}})$ are in agreement with the anomalous dimension presented in ref. [28]. The UV-sensitive logarithms for the quark case are presented here for the first time.

When electroweak UV logarithms are induced, the electroweak non-logarithmic terms in the matching at $\mu_{\text{ew}}$ that accompany them are scheme dependent and formally part of the next-to-leading-logarithmic (NLL) approximation. They should not be included in the analysis and we thus disregard them.

For those contributions that are not UV sensitive, we directly perform the matching of the SMEFT Lagrangian eq. (2.9) onto the effective Lagrangian eq. (3.2), integrating out the heavy degrees of the SM (top quark, Higgs, and the $W$ and $Z$ bosons). In all calculations we employ a general $R_\xi$ gauge for all gauge bosons and have verified the independence of our results of the gauge-fixing parameters. All calculations were performed using self-written FORM [29] routines, implementing when necessary the two-loop recursion presented in refs. [30, 31]. The amplitudes were generated using QGRAF [32] and FeynArts [33].

At tree level and leading order in the $1/\Lambda$ expansion, integrating out the Higgs induces the following non-zero initial conditions for four-quark operators (see first two diagrams in figure 1)

$$C_{qq}^q(\mu_{\text{ew}}) = \frac{m_q m_{q'}}{M_H^2} \frac{v^2}{\sqrt{2} m_q} \frac{v^2}{\Lambda^2} C_{q'H-}, \quad C_{q}^4(\mu_{\text{ew}}) = \frac{m_q^2}{M_H^2} \frac{v}{\sqrt{2} m_q} \frac{v^2}{\Lambda^2} C_{qH-},$$

with $q, q' = u, d, s, c, b$. In principle, also the analogous operators with leptons are induced ($\bar{q}q\ell\ell, \ell\ell\ell\ell$). However, since we will rely on a fixed-order computation to predict $d_e \propto C_{3e}^e$ (see below), these operators do not enter our analysis.

The matching at the electroweak scale also induces the Weinberg and dipole operators, but only at the two-loop level. The Weinberg operator obtains a two-loop initial condition
at $\mu_{\text{ew}}$ only from the top-quark operator (see third diagram in figure 1):

$$C_\ell(\mu_{\text{ew}}) = \frac{\alpha^2}{(4\pi)^2} \frac{v}{\sqrt{2m_f}} \frac{v^2}{\Lambda^2} C_{fH} - \frac{x_{th}}{2(1 - 4x_{th})} \left\{ 3 - 14x_{th} + 8x_{th}^2 \right\}
+ 6x_{th}(1 - 2x_{th} + 2x_{th}^2)\Phi \left( \frac{1}{4x_{th}} \right) + (2 + 10x_{th} - 12x_{th}^2) \log x_{th} \right\}, 
(4.4)$$

with $\Phi(z)$ defined in eq. (4.19). This result agrees with ref. [6], but is, as far as we are aware, presented here for the first time in a non-parametric form.

The electron photon dipole ($O^{\ell}_{3}$), and light-quark photon and gluon dipoles ($O^{q}_{3}, O^{q}_{4}$ with $q = u, d, s$) all receive two-loop initial conditions by integrating out the top quark and the Higgs, $W$, and $Z$ bosons:

$$C_3^{\ell}(\mu_{\text{ew}}) = \frac{\alpha^2}{(4\pi)^2} \sum_{f' = \ell, \mu, \tau, b, c, t} \frac{v}{\sqrt{2m_{f'}}} \frac{v^2}{\Lambda^2} C_{f'H} - A^{\ell}[f'],
(4.5)$$

$$C_3^{q}(\mu_{\text{ew}}) = \frac{\alpha\alpha_s}{(4\pi)^2} \sum_{f' = q, \tau, t} \frac{v}{\sqrt{2m_{f'}}} \frac{v^2}{\Lambda^2} C_{f'H} - A^{q}[f'],
(4.6)$$

$$C_4^{q}(\mu_{\text{ew}}) = \left( \sum_{q' = q, t} \frac{\alpha_s^2}{(4\pi)^2} B_{q'}^{q}[q'] + \frac{\alpha\alpha_s}{(4\pi)^2} B_{q'}^{q}[q] \right) \frac{v}{\sqrt{2m_{q'}}} \frac{v^2}{\Lambda^2} C_{q'H} - ,
(4.7)$$

with the precise decomposition of the coefficient functions $A^{\ell'}, B_{s}^{q}$, and $B_{ew}^{q}$ discussed below.

A few clarifications regarding eqs. (4.5)–(4.7) are in order. We calculate the electroweak matching in fixed-order perturbation theory whenever there are no QCD corrections that are enhanced by large logarithms related to light fermion masses. This is the case for all diagrams with top-quark loops, and the diagrams with lepton loops contributing to the electron dipole, i.e. the contributions $C_3^{\ell}(\mu_{\text{ew}})$ that are proportional to $C_{fH}$ - . (We also neglect the running of the electromagnetic coupling constant.) Therefore, the coefficients $C_{3/4}^{q}(\mu_{\text{ew}})$ do not receive threshold contributions from virtual quarks other than the top quark, since those are properly included via the RG-mixing of the four-quark operators (eq. (4.3)) into $C_{3/4}^{q}$, see below. In principle, also the contributions to $C_3^{\ell}(\mu_{\text{ew}})$ of all quarks other than the top should be obtained from the RG evolution; however, the corresponding

Figure 1. Sample Feynman diagrams contributing to the tree-level matching of four-quark operators (first two diagrams) and the two-loop matching inducing the Weinberg operator (third diagram). The red square vertex indicates a CP-violating Higgs Yukawa, $C_{fH}$.
mixed QED-QCD RG evolution is currently unknown.\(^2\) Thus, we temporarily include the contribution of bottom and charm quarks to \(C_3^\ell(\mu_{\text{ew}})\) with a fixed-order calculation, with the understanding that corrections to these results may be large. On the other hand, we do not include the corresponding fixed-order results for the light-quark \((u, d, s)\) contributions to \(C_3^\ell(\mu_{\text{ew}})\). In addition to the issue of the large unknown QCD corrections, these contributions are suppressed by the light-quark masses. Therefore, it is not \(C_3^\ell\) but the \(C_{3/4}^\ell\) coefficients that provide the numerically leading constraints on the \(C_{qH-}\) couplings for light quarks. Similarly, we do not include the contributions of the electron and muon to \(C_3^\ell(\mu_{\text{ew}})\). The appearance of their masses (that are much smaller than the lowest scale, \(\mu_{\text{had}}\), where these operators can be meaningfully defined) in the logarithms would spuriously enhance the corresponding bounds by large factors. These contributions should be taken into account properly by using the RG. The tau, on the other hand, has a mass large enough to justify a fixed-order calculation as an estimate and is thus included. Note, however, that numerically the strongest bounds on the coefficients of all three leptons would far arise from the electron EDM, and their contributions to the hadronic EDMs do not play a role in our analysis, given current experimental data.

The decomposition of the coefficients \(A^f\), \(B^q\), and \(B_{\text{ew}}^q\) in eqs. (4.5)–(4.7) are

\[
A^f(f') = \begin{cases} 
4n_cQ_f^2L_{1f} & \text{for } f = f' \\
4n_cQ_fL_{1f}^{-1}v_Zv_L & \text{for } f \neq f' 
\end{cases} \tag{4.8}
\]

\[
B^q[q] = \begin{cases} 
2L_{1q} & \text{for } q = q' \\
2L_{2q} & \text{for } q \neq q' = t 
\end{cases} \tag{4.9}
\]

\[
B_{\text{ew}}^q[q] = \left(\frac{v_Z^2 - a_Z^2}{s_Z}L_5 + \frac{v_Z^2 + a_Z^2}{s_Z}L_6\right) + \frac{1}{s_Z^2}L_7 \tag{4.10}
\]

with \(n_c[t] = n_c[q] = n_c\) and \(n_c[f] = 1\). The vector and axial-vector coupling of the \(Z\) boson to a fermion are \(v_f^2 = \frac{1}{2n_c} (T^f_1 - 2Q_fs_w^2)\) and \(a_f^2 = \frac{1}{2n_c} T^f_3\), respectively. In figure 2 and figure 3 we show representative diagrams contributing to \(L_1\)–\(L_7\). As discussed, only \(A^\ell[e]\) and \(A^q[q]\) contain UV sensitive electroweak contributions proportional to \(\log M^2_\ell/\Lambda^2\), thus all other scheme-dependent finite electroweak threshold corrections have been dropped in \(A^\ell[e]\) and \(A^q[q]\). We have, however, kept the terms from top-quarks loops (proportional to \(L_1\) and \(L_2\)) as they depend on the top-quark Yukawa and as such are independent of the

\(^2\) The QCD corrections to the bottom and charm contributions are naively estimated to be large (factor of a few in \(C_3^\ell\)) but require a more complex calculation [19].
Figure 2. Representative Barr-Zee-type Feynman diagrams contributing to the two-loop matching of dipole operators at the electroweak scale. The label under the diagrams indicates the loop function ($L_1$ and $L_2$) to which the corresponding diagram contributes. The red square vertex indicates a CP-violating Higgs Yukawa, $C_{fH}^-$. Electroweak diagrams like those in the second row are responsible for the UV sensitivity of contributions proportional to $C_{eH}^-$ and $C_{qH}^-$ for the electron and light-quark EDM operators, respectively. Diagrams with dimension-five couplings (lower right) are required to obtain a gauge-invariant result. See main text for details.

scheme-dependent electroweak corrections. The loop functions ($L_i$) read

\begin{align}
L_1 &= 2x_{th}(2 + \log x_{th}) + x_{th}(1 - 2x_{th})\Phi\left(\frac{1}{4x_{th}}\right), \\
L_2 &= \frac{x_lZx_{th}}{x_lZ - x_{th}} \left[(1 - 2x_{th})\Phi\left(\frac{1}{4x_{th}}\right) - 2\log x_{hZ} - (1 - 2x_{lZ})\Phi\left(\frac{1}{4x_{lZ}}\right)\right], \\
L_3 &= x_{f'h}\Phi\left(\frac{1}{4x_{f'h}}\right)^{x_f'h \ll 1} x_{f'h}\left(\log^2 x_{f'h} + \pi^2/3\right), \\
L_4 &= \frac{x_{f'\gamma}x_{f'h}}{x_{f'\gamma} - x_{f'h}} \left[\Phi\left(\frac{1}{4x_{f'\gamma}}\right) - \Phi\left(\frac{1}{4x_{f'h}}\right)\right] \\
&\quad \frac{x_f'Zx_f'h}{x_{f'\gamma} - x_{f'h}} (\log^2 x_{f'h} - \log^2 x_{f'\gamma}), \\
L_5 &= \frac{1}{3} \left[(3x_{hZ}^3 - 18x_{hZ}^2 + 24x_{hZ})\Phi\left(\frac{x_{hZ}}{4}\right) + 6x_{hZ}(1 + 2\log x_{hZ}) \\
&\quad + (12x_{hZ}^3 - 48x_{hZ}^2 + 12x_{hZ}) Li_2(1 - x_{hZ}) \\
&\quad + (x_{hZ}^3 - 4x_{hZ}^2)\pi^2 + (3x_{hZ}^3 - 12x_{hZ}^2)\log^2 x_{hZ}\right].
\end{align}
Figure 3. Representative Barr-Zee-type Feynman diagrams contributing to the two-loop matching of dipole operators at the electroweak scale. The label under the diagrams indicates the loop function ($L_3$–$L_7$) to which the corresponding diagram contributes. The red square vertex indicates a CP-violating Higgs Yukawa, $C_{fH}$. Here, $f$ denotes fermions and $q'$ quarks. The electroweak contributions to the chromo EDM operator from diagrams like the ones in the second row are UV finite as opposed to the ones in figure 2. See main text for details.

\[ L_6 = \frac{2}{9} x_Z^2 \left[ (3x_h^3 - 6x_h^2Z - 24x_hZ) \Phi \left( \frac{x_hZ}{4} \right) \right. \
+ (6x_h^2 + 24x_hZ) \log x_hZ + (6x_h^2Z - 24x_hZ) \right. \]
\[ + (6x_h^3 - 18x_hZ - 24) \text{Li}_2(1 - x_hZ) + (3x_hZ + 4)\pi^2 \right], \tag{4.17} \]

\[ L_7 = \frac{1}{18} x_W \left[ (3 - 6x_W - 24x_W^2) \Phi \left( \frac{x_W}{4} \right) + (3x_W^2 + 4x_W^3)\pi^2 \right. \
- (6x_W + 24x_W^2) \log x_W + (6x_W - 24x_W^2) \right. \]
\[ + (12x_W^3 + 9x_W^2 - 3)(2 \text{Li}_2(1 - x_W) + \log^2 x_W) \right], \tag{4.18} \]

where we defined the mass ratios $x_{ij} \equiv M_i^2 / M_j^2$. For the case of small fermion masses, i.e., when $x \ll 1$, we have expanded for convenience the $\Phi(z)$ function. Our loop functions $L_5$, $L_6$, $L_7$ correct a global factor of $\sqrt{2}$ with respect to the corresponding ones in ref. [8] and a typographical sign mistake in the $x_h^3$ coefficient of $\text{Li}_2(1 - x_hZ)$ in $L_5$ in the same reference. The loop function $L_2$ implicitly corrects an error in eq. (A.5) of ref. [28], where two logarithms seem to have been incorrectly added.
The function $\Phi(z)$ in the results above is given by \([30]\)

$$
\Phi(z) = \begin{cases} 
4 \sqrt{\frac{1}{1-z}} \text{Cl}_2(2 \arcsin(\sqrt{z})) & \text{for } 0 \leq z < 1, \\
\frac{1}{\lambda} \left(-4 \text{Li}_2(1-\frac{1}{\lambda^2}) + 2 \log^2(1-\frac{1}{\lambda^2}) - \log^2(4\lambda) + \frac{\pi^2}{3}\right) & \text{for } z > 1,
\end{cases}
$$

(4.19)

where $\lambda \equiv \sqrt{1-1/z}$, and the dilogarithm and Clausen function are defined as

$$
\text{Li}_2(x) = -\int_0^x du \log(1-u)/u \quad \text{and} \quad \text{Cl}_2(\theta) = -\int_0^\theta dx \log|2\sin(x/2)|,
$$

(4.20)

respectively.

Having computed the initial conditions at the electroweak scale, we perform the QCD RGE evolution from $\mu_{\text{ew}}$ to the hadronic scale $\mu_{\text{had}} = 2 \text{ GeV}$ (see refs. [6, 18] for details). For operators with quarks the RGE evolution resums large QCD logarithms and accounts for mixing among different operators. In the case of light-quark SMEFT operators, the tree-level induced four-quark operators ($O_{qq}^q$ and $O_{q1}^q$) mix at one-loop under QCD into the quark dipoles. Nevertheless, the contributions to $C_q^3(\mu_{\text{ew}})$ and $C_q^4(\mu_{\text{ew}})$ of two-loop diagrams with top-quark and $Z$-boson loops provide the numerically dominant effect \([23]\) as the four-fermion operators are additionally suppressed by an extra light-quark mass (see eq. (4.3)).

The situation is different for contributions from bottom- and charm-quark SMEFT operators. Here the leading contribution to partonic dipoles ($d_q, \bar{d}_q$) and the Weinberg operator ($w$) are induced by mixing during the RG evolution. The main reason is that the nuclear matrix elements of bottom and charm dipole operators are tiny. Therefore, diagrams like the ones in figure 2 do not contribute, i.e., bottom and charm quarks only enter as virtual particles in loop diagrams. At the same time, their mass is significantly below the electroweak scale. For this reason, a tree-level matching at the electroweak scale and the subsequent one-loop RG evolution \([27]\), which must include the mixing of four-fermion operators into dipole and Weinberg operators, is sufficient (and necessary) to obtain the leading-logarithmic result. The two-loop calculation has been performed in ref. [18] but is not used in this work, as perturbative uncertainties and higher-order corrections are generally neglected in our analysis (see the discussion in section 7.3).

5 Electric dipole moments

Non-zero coefficients of the higher dimension operators in eq. (2.9) induce in general electric dipole moments in nucleons, atoms, and molecular systems via contributions to the hadronic Lagrangian in eq. (3.1). Here we summarize the status of the induced dipole moments for the systems used in our fit, namely, the electron, neutron, and mercury EDMs.

In addition to these EDMs there are also constraints from the experimental measurement of the muon EDM \([34]\), as well as other systems with a hadronic component. The direct constraint from the current muon EDM measurement leads to constraints that are about six orders of magnitude weaker than the one obtained via virtual muon contributions to the electron EDM, and will thus not be used. Concerning other hadronic EDMs,
we have checked that within our setting the experimental bounds on the radium [35] and xenon [36] EDMs are not competitive with the neutron and mercury bounds and we do not include them in the fit either.

5.1 Electron

The most recent experimental bound on the electron EDM is [37]

\[ |d_e| < 1.1 \times 10^{-29} \text{ cm @ 90\% confidence level (CL)}. \] (5.1)

This value was obtained using ThO molecules, neglecting any CP-violating electron-nucleon couplings. Bounds on the electron EDM have also been obtained using YbF [38] and HfF$^+$ [39]. The resulting bounds are currently not competitive with the ThO bound and are thus not used in our fit.

5.2 Neutron

The simplest hadronic system used in our fit is the neutron. The most recent experimental bound on the neutron EDM is [40]

\[ |d_n| < 1.8 \times 10^{-26} \text{ cm @ 90\% CL}. \] (5.2)

The future projections estimate an improvement of the limit to $|d_n| < 10^{-27} \text{ cm}$ [40]. Throughout this work, we assume that the $\theta_{\text{QCD}}$ term has negligible effect on any EDMs and all effects arise from the Yukawa SMEFT operators. The dipole moments of the partons then contribute to the neutron EDM as

\[ \frac{d_n}{e} = (1.1 \pm 0.55)\langle \bar{d}_d + 0.5\bar{u}_u \rangle + \left( g_u^T \frac{d_u}{e} + g_d^T \frac{d_d}{e} + g_s^T \frac{d_s}{e} \right) \pm \frac{74}{3}(1 \pm 0.5)w \text{ MeV}. \] (5.3)

The hadronic matrix elements of the chromoelectric dipole operators ($\bar{d}_q$) and the Weinberg operator ($w$) are estimated using QCD sum rules and chiral techniques [24, 25, 41]. For the matrix elements of the electric dipole operators ($d_q$) we use the lattice results [42] $g_u^T = -0.204(15)$, $g_d^T = 0.784(30)$, $g_s^T = -0.0027(16)$. The sign of the hadronic matrix element of the Weinberg operator is not known; to be definite, we choose the positive sign in our analysis. Note that in our scenario the contribution of the Weinberg operator is always subdominant (either suppressed by small quark masses, or numerically small compared to the electron EDM bound), such that the sign ambiguity does not affect the results of our global fit.

5.3 Mercury

Significant constraints in our fit arise also from the mercury EDM. The experimental bound is [43]

\[ |d_{\text{Hg}}| < 7.4 \times 10^{-30} \text{ cm @ 95\% CL}. \] (5.4)

The relation of the partonic EDMs to that of mercury is given by [24]

\[ \frac{d_{\text{Hg}}}{e} = \kappa_S \left[ g_{\pi NN} \left( \frac{a_0}{e} - \frac{a_{1 T}}{e} \right) + \frac{d_n}{e} + \frac{d_p}{e} \right] + a_e \frac{d_e}{e}. \] (5.5)
Here, $\kappa_S = -2.8 \times 10^{-4} \text{fm}^{-2}$ denotes the contribution of the Schiff moment to the mercury EDM, with an error not exceeding 20% [44]. The expression in square brackets is the Schiff moment for mercury. The CP-odd isoscalar and isovector pion-nucleon interactions are given by $g_{\pi NN}^{(0)} = (5 \pm 10) \times (\bar{d}_u + \bar{d}_d) \text{fm}^{-1}$, $g_{\pi NN}^{(1)} = 20^{+40}_{-10} \times (\bar{d}_u - \bar{d}_d) \text{fm}^{-1}$, obtained from QCD sum-rule estimates [45], while $g_{\pi NN} \simeq 13.5$ [46, 47] is the CP-even pion-nucleon coupling. The contribution of these interactions to the Schiff moment is given by the parameters

$$a_0 = 0.01 \, e \, \text{fm}^3$$
and

$$a_1 = \pm 0.02 \, e \, \text{fm}^3.$$ We took these “best values” from ref. [24]; they have an intrinsic uncertainty of about an order of magnitude. The contributions of unpaired proton and neutron to the Schiff moment has been calculated in ref. [48], with result $s_p = 0.20(2) \, \text{fm}^2$ and $s_n = 1.895(35) \, \text{fm}^2$. Finally, the contribution of the electron EDM is subdominant; no explicit uncertainty is given for the prefactor $a_e = 10^{-2}$ [24, 45]; however, different evaluations lead to different signs [49]. See ref. [24] for a detailed discussion of all contributions to the mercury EDM. Employing “central values” for all parameters, we find numerically

$$\frac{d_{\text{Hg}}}{e} = -3.8 \times 10^{-4} \left[ 0.5(\bar{d}_u + \bar{d}_d) \pm 4(\bar{d}_u - \bar{d}_d) \right] - 5.3 \times 10^{-4} \frac{d_n}{e},$$

(5.6)

with $d_n/e$ given in eq. (5.3).

We have neglected several contributions in eq. (5.5). The isotensor pion-nucleon coupling contributes in principle, but is expected to be small in comparison to the scalar and vector coupling, as it arises at higher order in chiral perturbation theory [24]. Contributions of four-fermion operators are smaller than the chromo EDM contributions by two orders of magnitude, and are also neglected. Finally, in our analysis we neglected the proton EDM contribution that is one order of magnitude smaller than the neutron EDM contribution, as well as the small electron EDM contribution.

### 6 Collider observables

The SMEFT couplings in the Lagrangian in eq. (2.9) do not only induce electric dipole moments, but also affect the production cross sections and decay branching ratios of the Higgs boson. In this section we give an overview of the most important effects; the actual numerical implementation of the LHC constraints is discussed in section 7. The main effects are captured by considering modifications of the gluon fusion production cross section ($gg \to h$) and the $h \to \gamma\gamma$ branching ratio (both are one-loop induced both in the SM and in our setup), as well as of the branching ratios to fermions observed by ATLAS and CMS ($h \to bb$, $h \to cc$, $h \to \tau\tau$, $h \to \mu\mu$), and the total decay width of the Higgs. The strongest constraints on the $\tau$ Wilson coefficients are obtained from the angular distribution analysis of the $h \to \tau\tau$ measurement by CMS [50]. It is convenient to parameterise these modifications in terms of the parameters

$$r_{f,\pm} \equiv \frac{v}{\sqrt{2}m_f} \frac{v^2}{\Lambda^2} C_{fH \pm}.$$ (6.1)

Using eq. (2.11) we can readily translate $r_{f,\pm}$ to the $\kappa$-framework parameters, i.e.,

$$\kappa_f \cos \phi_f = 1 - r_{f,+}$$

and

$$\kappa_f \sin \phi_f = -r_{f,-}.$$
The deviation from the SM gluon fusion cross section or the decay to gluon-induced light jets can be effectively captured by (see ref. [6] for details)

\[ R_{gg} \equiv \frac{\sigma(gg \to h)}{\sigma(gg \to h)_{SM}} = \left( \left| \sum_q (1 - r_q) A(\tau_q) \right|^2 + \frac{3}{2} \sum_q r_q B(\tau_q) \right)\left/ \left| \sum_i A(\tau_i) \right|^2 \right. , \] (6.2)

where we have defined \( \tau_q \equiv \frac{4 m_q^2}{M_h^2} - i \varepsilon \), the sum runs over all quark flavours \( q = u, d, s, c, b, t \), and the fermionic one-loop functions are given by

\[ A(\tau) = \frac{3 \tau}{2} \left( 1 + (1 - \tau) \arctan^2 \frac{1}{\sqrt{\tau - 1}} \right) , \quad B(\tau) = \tau \arctan^2 \frac{1}{\sqrt{\tau - 1}} . \] (6.3)

Similarly, the modification of the Higgs decays into two photons with respect to the SM reads

\[ R_{\gamma\gamma} \equiv \frac{\Gamma(h \to \gamma\gamma)}{\Gamma(h \to \gamma\gamma)_{SM}} = \frac{\left| A_W + \frac{1}{6} \sum_f Q_f^2 n_c(f)(1 - r_f) A(\tau_f) \right|^2 + \left| \frac{1}{2} \sum_f Q_f^2 n_c(f) r_f B(\tau_f) \right|^2}{\left| A_W + \frac{1}{6} \sum_f Q_f^2 n_c(f) A(\tau_f) \right|^2} , \] (6.4)

where the sums run over all charged fermions \( f \) with \( n_c(f) = 3 \) for quarks and \( n_c(f) = 1 \) for leptons, and the bosonic contribution is given by

\[ A_W = -\frac{1}{8} \left( 2 + 3 \tau_W + 3 \tau_W (2 - \tau_W) \arctan^2 \frac{1}{\sqrt{\tau_W - 1}} \right) , \] (6.5)

with \( \tau_W = \frac{4 M_{W^2}}{M_{h^2}^2} - i \varepsilon \). We keep the modifications of the Higgs coupling to \( WW \) and \( ZZ \) bosons unmodified as the contributions from Yukawa operators are loop induced compared to the tree-level contributions already present at the dimension-four level. Moreover, \( h \to Z\gamma \) is omitted as its decay width is suppressed by a smaller coupling compared to \( h \to \gamma\gamma \).

The signal strengths of searches for Higgs decays to fermions are also affected by the corresponding modifications of the partial widths, namely:

\[ R_{ff} \equiv \frac{\Gamma(h \to f\bar{f})}{\Gamma(h \to f\bar{f})_{SM}} = (1 - r_{f, +})^2 + r_{f, -}^2 . \] (6.6)

As of now, the ATLAS and CMS collaborations have observed Higgs decays into \( f = b, c, \tau, \mu \). However, the modifications of the partial widths to any fermion affect the total width of the Higgs via

\[ \frac{\Gamma_{tot}}{\Gamma_{tot, SM}} = 1 + \sum_f (R_{ff} - 1) \text{Br}(h \to f\bar{f})_{SM} + \sum_{X=gg,\gamma\gamma} (R_X - 1) \text{Br}(h \to X)_{SM} , \] (6.7)

which in turn affects all its branching ratios, i.e., all the signal strength for Higgs searches.

We conclude this section with some remarks concerning the convergence of the SMEFT expansion and the potential impact of dimension-eight operators. The issue arises because the squared amplitude enters the expression (6.6), leading to terms proportional to \( 1/\Lambda^4 \)
in $R_{ff}$ that we keep in our numerics. What would be the impact of including dimension-eight Yukawa operators (of the generic form $Q^{(8)} \sim (H^\dagger H)Q_L H C_{\mu H f R}$) in the amplitude? While this question is hard to answer quantitatively without actually performing the analysis, the following arguments suggest that the impact would result only in minimal changes on the bounds of the dimension-six coefficients.

First, note that this issue concerns only the real parts of the Wilson coefficients. As there is no interference with the SM in the CP-odd part of the amplitude, such dimension-eight terms would be of order $1/\Lambda^8$. The first term on the right side of eq. (6.6), on the other hand, would be changed to

$$(1 - r_{f,+})^2 \to 1 - 2C^{(6)} \frac{v}{\sqrt{2}m_f} \frac{v^2}{\Lambda^2} + (C^{(6)})^2 \frac{v^2}{2m_f^2} \frac{v^4}{\Lambda^4} - C^{(8)} \frac{v}{\sqrt{2}m_f} \frac{v^4}{\Lambda^4} + O(\Lambda^{-6}). \quad (6.8)$$

Here, $C^{(6)}$ and $C^{(8)}$ denote generic dimension-six and dimension-eight Wilson coefficients, respectively. We see that the additional last term in eq. (6.8) is parametrically suppressed with respect to the linear dimension-six term by a factor $v^2/\Lambda^2$, and with respect to the quadratic dimension-six term by a factor $m_f/v$. A potential problem arises if the term quadratic in $C^{(6)}$ dominates the fit. A direct comparison between the last two terms in eq. (6.8), under the “EFT assumption” $C^{(6)} \sim C^{(8)}$, shows that the dimension-eight contribution will be subleading as long as $C^{(6)} > \sqrt{2}m_f/v$. We will see in figure 4 that this condition is clearly satisfied for all fermions apart from the top quark. While this is only a rough estimate, we interpret this as an indication that the impact of the dimension-eight terms on the fit would be tiny. By contrast, figure 5 shows that for the top quark $C_{tH+} \lesssim 3$, which is larger than $\sqrt{2}m_t/v \sim 1$, such that for values of $C^{(8)} \sim 9$ the dimension-eight term would contribute significantly to the fit. While this required value of $C^{(8)}$ is close to the perturbativity limit, we conclude that the EFT expansion in the collider observables might not work as well for the top quark as it does for all other fermions. CP-odd contributions, on the other hand, are not expected to receive large corrections from dimension-eight operators for any of the fermions, as discussed above.

7 Global analyses

Having obtained in the previous sections the expressions for the relevant observables, EDMs, Higgs production cross sections ($\sigma$), and branching ratios (BR) decay widths we now combine the constraints on the considered Wilson coefficients in a global analysis. For this purpose we use the GAMBIT global fitting framework [15] to calculate a combined likelihood based on these two sets of constraints. The collider likelihoods are taken from the HiggsSignals_2.5.0 and HiggsBounds_5.8.0 codes [51, 52] interfaced with the ColliderBit module [16] of GAMBIT.

HiggsSignals contains three modules to compute likelihoods from Higgs measurement at the LHC. Each module calls HiggsBounds to calculate the various production cross sections and branching ratios within our SMEFT model and to obtain the signal strengths implemented in each module.
The first module computes a likelihood of Run 1 Higgs measurements using a set of signal strengths provided by the ATLAS and CMS combination of Run 1 data [53]. In this module, the signal strengths are “pure channels” in the sense that one decay channel is combined with one specific production channel, i.e.,

\[ \mu_i^f = \frac{\sigma_i \cdot BR_i^f}{\sigma_i \cdot BR_{i,SM}^f} \]

with \( i \) and \( f \) indicating different production modes (\( \text{ggh, VBF, Wh, Zh, tth} \)) and decay channels (\( \text{ZZ, WW, } \gamma \gamma \tau \tau, \bar{b}b \)), respectively. Run 1 data are sensitive to 20 such signal strengths. The (Log)likelihood is then obtained from

\[ \chi^2 = (\mu - \mu^{\exp})^T C_{\text{cov}}^{-1} (\mu - \mu^{\exp}), \quad (7.1) \]

where \( \mu \) are the vectors containing the 20 signal strengths computed as a function of the SMEFT Wilson coefficients, \( \mu^{\exp} \) are the corresponding experimental combinations, and the superscript “\( T \)” denotes transposition. The matrix \( C_{\text{cov}} \) is the signal-strength covariance matrix describing the uncertainties of and correlations among the signal strengths.

For Run 2 measurements there are two HiggsSignals modules to compute a likelihood, one that uses the simplified template cross section measurements (STXS) [54] and another that uses the peak-centered method. In our analysis we mainly use the former since the peak-centered method increases the computing time by almost an order of magnitude and has also less constraining power, as we have checked explicitly. The only exception in which we do use the peak-centered-method module is to include the two \( h \rightarrow \mu\mu \) analyses containing in total 34 peak measurements [55, 56]. Each signal strength is a combination of various production modes weighted by the corresponding experimental efficiency (\( \epsilon_i \)), which accounts for the detector performance in identifying signal events, i.e.,

\[ \mu = \frac{\sum_j N_j (\sigma \cdot BR)_j \epsilon_j}{\sum_j (\sigma_{\text{SM}} \cdot BR_{\text{SM}})_j \epsilon_{\text{SM}}^j}. \quad (7.2) \]

In our analysis we include the signal strengths of 56 measurements originating from experimental searches from June 2021 [57–76]. The corresponding (Log)likelihood is then obtained in an analogous manner as for the Run 1 data in eq. (7.1). The SM predictions for the cross sections are obtained through a fit to the predictions from Yellow Report 4 [54].

The analysis using the angular correlation between the decay planes of two taus from Higgs decays [50] is not included in HiggsSignals; we have implemented it separately. As this angular analysis is not strongly correlated with any of the other branching ratio measurements, we simply add this one \( \chi^2 \) value.

To compute a likelihood from the EDM measurements, we assume their experimental uncertainties \( \sigma^{\exp} \) to be Gaussian distributed yielding the likelihoods (log \( \mathcal{L} \))

\[ \chi^2_X = -2 \log \mathcal{L}_X = \frac{(d_X(2 \text{ GeV}) - d_X^{\exp})^2}{\sigma_X^{\exp}^2}, \quad (7.3) \]

with \( X = n, e, \text{Hg} \). The uncertainties \( \sigma_{X,\text{exp}} \) are obtained from the upper limits given in eqs. (5.1), (5.2), and (5.4) by assuming zero central values for \( d_X^{\exp} \).

The total likelihood of a parameter point is computed from the sum of the \( \chi^2 \) values from the EDM and LHC likelihoods. To identify preferred regions in a subspace of the
scanned parameters we must profile over the remaining scanned parameters. We perform the profiling using \texttt{pipipi} \cite{77}, which computes the lowest possible $\chi^2$ value of each parameter point in the subspace of interest by allowing all other parameters to float simultaneously to minimize the $\chi^2$. When projecting onto a two-parameter subspace, the allowed regions at 68% and 95% CL correspond to the parameter space for which the difference $\chi^2 - \chi^2_{\text{best}}$ is less or equal than $\chi^2_{68\%} \approx 2.28$ and $\chi^2_{95\%} \approx 5.99$, respectively. $\chi^2_{\text{best}}$ is the $\chi^2$ value of the best-fit point, i.e., the lowest $\chi^2$ value. In our case the $\chi^2$ value of the SM (all SMEFT Wilson coefficients are zero) divided by the number of the included $N_{\text{d.o.f.}} = 79$ observables is $\chi^2_{\text{SM}}/N_{\text{d.o.f.}} = 0.83$, where we did not include the 34 peak observables from the $h \to \mu\mu$ measurements and the angular analysis of $h \to \tau\tau$; they only affect the muon and tau Yukawas, respectively. In the scan that we present the difference between the $\chi^2$ of the best-fit point and the SM is always less than 0.1, so we do not display the best-fit value in the plots.

7.1 One-flavour scans

To set the stage, we first perform two-dimensional scans over the Wilson coefficients of each fermion individually, with the Wilson coefficients of all other fermions set to zero. The results are shown in figure 4. We allow the Wilson coefficients to float within their perturbative values, $|C_fH| \leq 4\pi$, and fix $\Lambda = 1\,\text{TeV}$ in our scan. All numerical input parameters are taken from ref. \cite{78}.

In this work, we parameterize the effects beyond the SM by the SMEFT Wilson coefficients in eq. (2.9). EDMs only constrain the imaginary part of the Wilson coefficients, $C_{fH-}$, while LHC observables constrain also the real parts, $C_{fH+}$. In general, the electron EDM gives the strongest bounds on CP-violating parts of the lepton and heavy-quark (top, bottom, charm) coefficients, while the corresponding light-quark (up, down, strange) coefficients are mainly constrained by a combination of neutron and mercury EDM measurements.

• Top

The constraints from the modification of Higgs production in gluon fusion have the parametric dependence $\mu_{gg} = (1 - r_t)^2 + r_t^2$, where we used the asymptotic values $A(\infty) = B(\infty) = 1$, valid to good approximation for the top quark, and neglected the contribution of all other quark flavours. In the same approximation, we have $\mu_{\gamma\gamma} = (1.0 + 0.27 r_{t,+})^2 + 0.17 r_{t,-}^2$. These are the dominant LHC constraints on the top couplings; note, however, that all production channels are included in our analysis. The corresponding constraints are shown in figure 4 (bottom left panel). The strongest individual bound on the CP-odd coefficient $C_{tH-}$ arises from the electron EDM, while the constraints arising from the hadronic systems are much weaker, see figure 4. A zoomed-in version of the combined fit region is shown in figure 5.

• Bottom, charm, tau, muon

None of the bottom, charm, tau, and muon EDMs themselves have been measured with high precision. Consequently, EDM bounds on CP-odd Higgs couplings of the fermions arise from their virtual contributions to the electron EDM (again, the bounds arising from
2D single-flavour scans \((\Lambda = 1 \text{ TeV})\)

Figure 4. EDM and LHC constraints on CP-even \((C_{fH^+}\)) and CP-odd \((C_{fH^-}\)) Higgs Yukawa couplings assuming \(\Lambda = 1 \text{ TeV}\). Each plot shows the result of a 2D scan in which only the two Wilson coefficients of the respective fermion are sampled. The contours represent the allowed 68\% and 95\% confidence regions for their respective observables. The colour coding of individual constraints is given in the legend, grey/black areas correspond to the regions allowed by the combination of all constraints. Regions allowed by an individual observable are only shown if the observable has a constraining effect on the parameter space. The sampled intervals are bounded by the perturbativity condition \(|C_{fH^{\pm}}| \leq 4\pi\).

hadronic systems are negligible in comparison). It is remarkable that the presence of a large quadratic logarithm (see eq. (4.13)) makes the electron EDM bound weaker than that on the top by only a factor 5 for the bottom, a factor 2 for the charm, a factor 6 for the
muon, and of the same order for the tau, rather than being suppressed by the much larger factor $Q^2 m_t/Q_t^2 m_f = \mathcal{O}(50 - 500)$, expected naively [28] from the Yukawa suppression. For the bottom and charm quarks, this strong logarithmic enhancement indicates that QCD corrections are large in these cases and should be included [19], similar to the case of hadronic dipole moments [6, 18] (see the discussion in section 3).

Measurements at the LHC directly constrain the decay of the Higgs into bottom, charm, tau, and muon pairs. The main effect comes from the modification of the partial $h \to f \bar{f}$ widths (see eq. (6.6)). This is a simplified picture, as modifying any Yukawa also affects the total Higgs decay width (see eq. (6.7)), the $h \to \gamma\gamma$ decay and, for quarks, also Higgs production via gluon-fusion. This effect is largest for the bottom quark (that dominates the SM Higgs decay width); nevertheless, we include the effect of all fermions on the total Higgs width.

Note that a negative value for the bottom Yukawa (in the sense of eq. (2.11)) is excluded at 68% CL (bottom middle panel in figure 4). Note also that for the muon Yukawa, the recent LHC measurements are more constraining than the electron EDM by an order of magnitude [79]. Regarding tau couplings, the CMS analysis on the CP structure of the $h \to \tau\tau$ decay [50] disfavours large values of $|C_{\tau H-}|$, and the 2$\sigma$ LHC constraint (blue region in lower right panel of figure 4) is split into two distinct regions. However, the constraint on $|C_{\tau H-}|$ from the electron EDM is still stronger.

- **Electron, up, down, strange**

The main EDM bounds on the electron and light-quark (up, down, strange) SMEFT coefficients arise from their contributions to the electron EDM and the neutron and mercury EDMs, respectively, as discussed in section 4. Note the different impact of the neutron and mercury EDMs on the up and down coefficients, due to the strong isospin dependence of the corresponding EDM predictions.
LHC bounds on the Wilson coefficients for the electron and the up and down quarks arise from modifications of the total Higgs decay width eq. (6.7). Indeed, we checked that the contributions to gluon fusion and $h \rightarrow \gamma\gamma$ are subleading. (For modifications of Higgs production induced by parton distribution functions, see ref. [80].) However, the resulting constraints are very weak when compared to the SM Yukawas, as illustrated by the corresponding ratios $|r_{e,+}| \lesssim 4000$, $|r_{u,+}| \lesssim 500$, and $|r_{d,+}| \lesssim 225$. In contrast, the bound on the CP-odd strange-quark coefficient from LHC is almost competitive with the hadronic EDM bounds (central panel in figure 4).

7.2 Two-, three-flavour scans, and beyond

Next, we let the Wilson coefficients of more than one fermion flavour float simultaneously. This allows for the cancellation of the contributions of the considered Wilson coefficients to the constraining EDMs, thereby relaxing the bounds in certain regions of parameter space. There are no such cancellations possible in the collider observables, because the main effect comes from partial decay widths where no interference is possible. (The small interference term between top- and bottom-quark contribution to Higgs production via gluon fusion and $h \rightarrow \gamma\gamma$ is negligibly small.) However, the bounds on the different Wilson coefficients are still correlated. For instance, the Higgs decay into bottom quark dominates the total Higgs decay width, and thus the $h \rightarrow b\bar{b}$ rate affects all branching ratios significantly. The same is true, albeit to a lesser extent, for all other Higgs decays.

• Up and down
In the left panel of figure 6 we show the results of a two-parameter scan over $C_{uH^-}$ and $C_{dH^-}$. We do not scan over the corresponding CP-even Wilson coefficients as they do not enter the EDM predictions. The different dependence of the neutron and the mercury EDMs on the up- and down-quark coefficients is clearly visible. With the isoscalar pion-nucleon coupling (entering the mercury EDM prediction) being subdominant, the bands of the two observables are nearly orthogonal, thus allowing to set stringent constraints on both $C_{uH^-}$ and $C_{dH^-}$.

• Bottom and strange
In the right panel of figure 6 we show the results of a four-parameter scan over $C_{bH^\pm}$ and $C_{sH^\pm}$ after profiling over the CP-even couplings. As we neglect the tiny strange-quark contribution to the electron EDM, the only constraint on $C_{sH^-}$ arises from the two hadronic EDMs. On the other hand, $C_{bH^-}$ receives its dominant constraint from the electron EDM, while the contributions to the hadronic EDMs are also taken into account.

We include the CP-even Wilson coefficients in the scan as they are both bounded by LHC measurements. As there is no direct measurement of $h \rightarrow s\bar{s}$, the correlation between $C_{sH^-}$ and $C_{bH^-}$ results from the contributions to the total Higgs decay width. The analogous plots with CP-even Wilson coefficients do not contain any additional information compared to the one-flavour scans and are thus not shown here.
• Top and bottom; top and tau; top and charm
In figure 7 we show the results of three different four-parameter scans, floating $C_{tH^\pm}$ simultaneously with $C_{bH^\pm}$ (first column), with $C_{\tau H^\pm}$ (second column), and with $C_{cH^\pm}$ (third column). The electron EDM bounds on $C_{tH^-}$ could, in principle, be lifted by two orders of magnitude compared to the single-flavour scan by choosing values close to the perturbativity limit for $C_{bH^-}$, $C_{\tau H^-}$, and $C_{cH^-}$. However, given the LHC bounds on these parameters the bound on $C_{tH^-}$ is weakened by only a factor of the order of five. Note also that allowing $C_{tH^+}$ to float significantly increases the allowed range of values for $C_{bH^+}$ at the 68% CL (upper left panel in figure 7). The relaxation of the bounds compared to the single-flavour fit is even more pronounced for the charm quark (upper right panel in figure 7). In the last row, we present the constraints on the parameters $C_{tH^+}$ and $C_{tH^-}$, with the other parameters profiled. This can be directly compared to figure 5, showing a significant relaxation of the bounds.

Interestingly, the combination of the electron EDM with LHC constraints has a profound impact on the constraint on the CP-even Wilson coefficients and not only on the CP-odd ones as one would naively expect. To understand this better, we first consider the second panel from above in the first column of figure 7 showing the allowed $C_{tH^-}$–$C_{bH^-}$ region. Here, the allowed combined region corresponds to the intersection of the electron EDM and the LHC constraint, which implies that for each allowed pair of $(C_{tH^-}, C_{bH^-})$, it is possible to find corresponding allowed values of $C_{bH^+}$ and $C_{tH^+}$ (which have been pro-
Figure 7. Constraints resulting from a 4D scan of top and bottom coefficients ($C_{tH^\pm}, C_{bH^\pm}$; first column), top and tau coefficients ($C_{tH^\pm}, C_{\tau H^\pm}$; second column), and top and charm coefficients ($C_{tH^\pm}, C_{cH^\pm}$; third column), assuming $\Lambda = 1$ TeV. In each plot only two parameters are shown, the remaining two are profiled over (see main text). Contours represent the allowed 68% and 95% confidence regions, the colour coding of individual constraints is given in the legend, and grey/black areas correspond to the combined regions. For details see main text and the caption of figure 4.
filed out in the plot). This can easily be verified by looking at the bottom left and bottom center panels of figure 4. By contrast, the electron EDM on its own does not constrain the $C_{tH^-} - C_{bH_+}$ subspace at all, i.e., the whole parameter space of the third panel in the first column of figure 7 ($C_{tH^-} - C_{bH_+}$ plot) is allowed with respect to the electron EDM. The reason being that one can always cancel the top against the bottom contributions to the EDM (see green band in the upper left panel). However, the combined EDM-LHC region is smaller than the one allowed by LHC alone, as not all values of $C_{bH^-}$ required to cancel the contributions of $C_{tH^-}$ are allowed by LHC bounds. In fact, the bottom single-flavour analysis (bottom center panel of figure 4) indicates that roughly $C_{bH^-} \approx C_{bH_+}$, resulting in the much smaller allowed combined region in the 4D scan. Similar arguments apply to the plots that show the top-tau and top-charm coefficients. Regarding the case of the charm quark, note that its contribution to the electron EDM is larger than that of the bottom quark, while the LHC constraint is comparatively weaker, resulting in a larger allowed combined region. Finally, we remark that while the contribution to the electron EDM of the muon is similar to that of the bottom, the muon Yukawa is so strongly constrained by recent LHC measurement that no appreciable cancellation can occur, which is why we do not present this scan.

- **Bottom and tau**
  In figure 8 we show the results of a scan of the four parameters $C_{bH_{\pm}}$ and $C_{\tau H_{\pm}}$. As in the previous four-parameter scans, there is an interesting interplay between EDM and LHC bounds. When considering EDM bounds only, we can always cancel the constraint if either $C_{bH^-}$ or $C_{\tau H^-}$ is profiled. Hence, there are no pure EDM constraints in any but the upper center panel where the CP-odd coefficients $C_{bH^-}$ and $C_{\tau H^-}$ are displayed. In contrast, the EDM constraints cannot always be satisfied if also LHC constraints are included. In the upper left plot with the two bottom coefficients displayed one can see that the allowed, combined region is enlarged compared to the single-flavour scan. However, for extreme values for $C_{bH^-}$, still allowed by LHC bounds, $C_{\tau H^-}$ cannot be profiled such that it compensates the bottom contribution to the electron EDM and simultaneously still be within the $2\sigma$ level $\tau$ LHC bounds. In contrast, $C_{bH^-}$ can be profiled such that the $C_{\tau H^-}$ coefficient in the upper right plot is only bounded by LHC constraints.

  Even though the tau contributions to the electron and quark EDMs are larger than those of the bottom by about a factor $3m_\tau \log^2(m_\tau/M_h)/m_b \log^2(m_b/M_h) \approx 2$ (apparent from eq. (4.5)), the allowed range for $C_{\tau H^-}$ is so strongly bounded by LHC constraints that the tau contribution can always be fully cancelled by the bottom contribution, but not vice versa. This implies that, given current data, the combined bounds (apart from the combination $(C_{\tau H^+}, C_{\tau H^-})$) are more stringent than either the LHC bounds or the EDM bounds alone. The effect of the electron EDM, further restricting the parameter regions allowed by LHC data, is also clearly visible in the combined region in the lower center panel that shows the bounds on the two CP-even coefficients $C_{bH_+}$ and $C_{\tau H_+}$.

- **Charm and tau**
  In figure 9 we present the results of a scan of the four parameters $C_{cH_{\pm}}$ and $C_{\tau H_{\pm}}$. The results are analogous to the case of bottom and tau discussed above. Note that the charm
4D $b, \tau$-scan \quad ($\Lambda = 1$ TeV)

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure8}
\caption{Constraints resulting from a 4D scan of bottom and $\tau$ Wilson coefficients ($C_{bH^\pm}$, $C_{\tau H^\pm}$) assuming $\Lambda = 1$ TeV. In each plot only two parameters are shown, the remaining two are profiled over (see main text). Contours represent the allowed 68\% and 95\% confidence regions, the colour coding of individual constraints is given in the legend, and grey/black areas correspond to the combined regions. For details see main text and the caption of figure 4.}
\end{figure}

The contribution to the electron EDM is larger than the bottom contribution by a factor of roughly $4m_c \log^2(m_c/M_h)/m_b \log^2(m_b/M_h) \approx 1.6$, while the LHC bounds on the charm quark are considerably weaker that those on the bottom quark.

• Top, bottom, and tau (third generation)

In figure 10 we present a scan of all six third-generation parameters $C_{bH^\pm}$, $C_{\tau H^\pm}$, and $C_{tH^\pm}$. It is interesting to compare the results to the case in which only two out of the three flavours were included in the fit (figures 7, 8).

First, we focus on the set of the four panels (middle and lower row, left and center) that show the same parameter combinations as the panels at similar positions in figure 7. We see that after profiling the remaining third-generation couplings, an “indirect” electron EDM constraint on $C_{tH^-}$ remains, although the allowed region is now significantly larger. By contrast, the “indirect” electron EDM constraint can be lifted completely by $C_{tH^-}$ (that is only weakly constrained from LHC measurements) in the panels in the top row and center right of figure 10, leaving only the LHC constraints. This should be compared to the corresponding much smaller combined regions in figure 8.
Figure 9. Constraints resulting from a 4D scan of charm and τ Wilson coefficients \( (C_{cH\pm}, C_{\tau H\pm}) \) assuming \( \Lambda = 1 \) TeV. In each plot only two parameters are shown, the remaining two are profiled over (see main text). Contours represent the allowed 68% and 95% confidence regions, the colour coding of individual constraints is given in the legend, and grey/black areas correspond to the combined regions. For details see main text and the caption of figure 4.

The bottom right panel of figure 10 shows the constraints on both top Wilson coefficients, \( C_{tH\pm} \). Notice that the allowed values increased by about a factor of two for \( C_{tH+} \) and by one order of magnitude for \( C_{tH-} \) compared to the single-flavour scan (figure 5).

The large computational resources required for scans with more than six parameters prevent us from scanning over more than three flavours. Nevertheless, we can infer some results in this direction from the one-, two-, and three-flavour scans above. Specifically, we will consider how the constraints on the third-generation Wilson coefficients change when including more flavours in the scans.

- **Top, bottom, (electron, light quark)**
  As an example, we float \( C_{tH\pm} \) and \( C_{bH\pm} \) under the presumption that any contributions to EDMs can be compensated by appropriate values of the coefficients \( C_{eH-} \) and \( C_{qH-} \), \( q = u, d \), which we, however, do not scan over. In this way only LHC bounds remain, which are one to two orders of magnitude weaker for \( C_{tH-} \) compared to the electron EDM. The results are shown in figure 11. We find that the LHC constraints on \( C_{tH+} \) and \( C_{bH+} \)
Figure 10. Constraints resulting from a 6D scan of top, bottom, and $\tau$ Wilson coefficients ($C_{tH^\pm}$, $C_{bH^\pm}$, $C_{\tau H^\pm}$) assuming $\Lambda = 1 \text{ TeV}$. In each plot only two parameters are shown, the remaining ones are profiled over (see main text). Contours represent the allowed 68% and 95% confidence regions, the colour coding of individual constraints is given in the legend, and grey/black areas correspond to the combined regions. For details see main text and the caption of figure 4.

are somewhat lifted compared to the single-flavour scans due to the large contribution of the bottom to the total Higgs width, while the absence of EDM bounds allows for much larger allowed ranges of $C_{tH^-}$ and $C_{bH^-}$. The analogous results for the charm instead of the bottom couplings are shown in figure 12.

More generally, the CP-violating up, down, and electron Wilson coefficients are merely and severely constrained by EDM measurements. Hence, including $C_{eH^-}$, $C_{uH^-}$, and $C_{dH^-}$ in a scan over the Wilson coefficients of the second or third generation would allow to completely cancel any EDM constraints and again only LHC constraints would remain.
Figure 11. LHC constraints resulting from a 4D scan of top and bottom Wilson coefficients ($C_{tH\pm}, C_{bH\pm}$) assuming $\Lambda = 1$ TeV. In each plot only two parameters are shown, the remaining ones are profiled over (see main text). Contours represent the allowed 68% and 95% confidence regions, the colour coding of individual constraints is given in the legend, and grey/black areas correspond to the combined regions. For details see main text and the caption of figure 4.

7.3 Theory uncertainties

In this work, we did not study the impact of theoretical uncertainties on the bounds on the Wilson coefficients. Hence, a short discussion of these effects is in order. The relevant uncertainties are: (i) uncertainties in the hadronic matrix elements; and (ii) perturbative uncertainties.

(i) The uncertainties on the hadronic matrix elements have been shown in section 5. Note that, in addition to the ranges given for the parameters, also some of the relative signs are not determined. In our numerical analysis, we have taken the central values for the hadronic matrix elements, and (somewhat arbitrarily) chosen the positive signs where they were not determined. We did not include these uncertainties in our likelihood function. In fact, no bounds at the 68% CL would result from the mercury EDM, while the neutron EDM bounds would get weaker by about a factor 2. At 95% CL, there would also be no bound from the neutron EDM. (In ref. [2], much smaller effects of the hadronic uncertainties have been found. This may be related to the statistical
Figure 12. LHC constraints resulting from a 4D scan of top and charm Wilson coefficients ($C_{tH^\pm}$, $C_{cH^\pm}$) assuming $\Lambda = 1$ TeV. In each plot only two parameters are shown, the remaining ones are profiled over (see main text). Contours represent the allowed 68% and 95% confidence regions, the colour coding of individual constraints is given in the legend, and grey/black areas correspond to the combined regions. For details see main text and the caption of figure 4.

(ii) It is important to recognize that, in many cases, the perturbative uncertainties are as large as the non-perturbative uncertainties. As an example, consider the bounds on the bottom and charm Wilson coefficients, studied (within the $\kappa$ framework) in ref. [18]. There it was shown that the QCD corrections are large; after inclusion of the two-loop leading-logarithmic QCD corrections, the uncertainties on the electric and chemolectric Wilson coefficients are reduced to order of 30%. We do not include the NLO corrections here, as lattice results are not available for all required matrix elements (see refs. [81–83] for preliminary results). Maybe somewhat surprisingly, also
the electron EDM bound on the CP-violating bottom and charm Yukawas receives large QCD corrections. The calculation of these effects is ongoing [19] and also not included in this work. No theory uncertainties are included in the LHC constraints. Note that they are partially contained in the uncertainties quoted by the experiments.

8 Discussion and conclusions

In this work we have presented the first high-dimensional fits of the coefficients of Yukawa-type SMEFT operators to multiple EDMs and LHC data. As expected, upon inclusion of a sufficient number of Wilson coefficients, all EDM constraints can be evaded by cancelations, and only LHC bounds remain. However, when considering the heavy fermions only, a nontrivial interplay between contributions remains.

There are several ways to further extend our analysis in the future. As discussed in section 7.3, we have neglected the impact of hadronic and perturbative uncertainties on our fit results. This impact can be profound in the lower-dimensional scans, and further motivates the ongoing efforts to decrease the uncertainties.

Improved experimental bounds (or discoveries), as well as the inclusion of additional EDMs such as those of proton and deuterium, once they become available, are expected to have a significant impact on the global fit. In particular, the different isospin dependence of the various hadronic systems will help to disentangle CP-odd Higgs couplings of the first-generation quarks. The CP structure of the heavy-fermion Yukawas can be tested more directly at present and future colliders by studying observables that are designed specifically to test the CP-odd couplings and typically require a vast amount of (expected) future data (see, e.g., refs. [84–90]).

Finally, as explained in section 2 where we introduced the theoretical framework for flavour-diagonal CP-violation in Higgs Yukawas, we do in general expect new flavour violating sources to accompany beyond-the-SM CP-violation. Therefore, allowing for flavour-changing contributions within SMEFT, and including the corresponding observables would be a further extension of the current analysis. While such contributions are generically expected in realistic UV models and are expected to lead to much stronger bounds, the question is whether, given the proliferation of parameters and observables, it is not better to study selected UV models directly.

In the final stages of this work, ref. [91] was published, presenting analyses with a similar scope as ours. We briefly comment on the differences between the two papers. In ref. [91], the κ framework (discussed in this paper at the end of section 2) is employed to perform multi-parameter fits to LHC data. The interplay of LHC data and the electron EDM bound is discussed, but no combined fit to both LHC and EDM data is performed. By contrast, constraints that arise from considerations of electroweak baryogenesis are discussed in ref. [91]. The expression for the electron EDM (eq. (13) in ref. [91]) is given in numerical form and thus hard to verify. However, it seems that the gauge-boson contribution must be incorrect (either gauge-dependent or containing an implicit logarithmic dependence on the UV cutoff that is not made explicit; the cited references contain contradictory results). See our discussion in section 4.
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A The $R_\xi$-gauge Lagrangian

In eq. (2.9) we presented the Yukawa Lagrangian that includes SMEFT modification from Yukawa operators in the mass-eigenstate basis and in unitarity gauge. In our computation we work in generic $R_\xi$ gauge. In what follows we present the relevant part of the Lagrangian, $\mathcal{L}_{\text{Yukawa}}$, in $R_\xi$ gauge after rotating to the mass-eigenstate basis for the fermions. We split the Lagrangian into

$$\mathcal{L}_{\text{Yukawa}} = \sum_{f=u,d,e} (\mathcal{L}_{\text{dim-4},f} + \mathcal{L}_{\text{dim-5},f} + \mathcal{L}_{\text{dim-6},f}),$$

(A.1)

where

$$\mathcal{L}_{\text{dim-4},u} = -m_u \bar{u} u - \bar{u} \left( \frac{m_u}{v} + \frac{v}{\sqrt{2}\Lambda^2} C u H^+ \right) u h$$

(A.2)

$$\mathcal{L}_{\text{dim-5},u} = \frac{v}{2\sqrt{2}\Lambda^2} \bar{u} \left( C_{u h}^+ + G_0^2 + 2G^+ G^- \right) u$$

(A.3)

$$\mathcal{L}_{\text{dim-6},u} = \frac{1}{2\sqrt{2}\Lambda^2} \left( C_{u h}^+ + G_0^2 + 2G^+ G^- \right) \bar{u} \left( C_{u h}^+ + G_0 \right) u$$

(A.4)

Here, $u = (u, c, t)$, $d = (d, s, b)$, $m_u = (m_u, m_c, m_t)$, and the $C_{uH}$ can also contain flavour off-diagonal pieces. The corresponding Lagrangian for down-type quarks is obtained by the obvious interchanges $u \leftrightarrow d$, $G^+ \leftrightarrow G^-$, $V^\dagger \leftrightarrow V$, and by a reversal of sign in all terms that are odd in the Goldstone fields. Analogously, the corresponding Lagrangian for leptons is obtained from the up-type Lagrangian above by the obvious interchanges $u \rightarrow e$, $d \rightarrow \nu$, $G^+ \leftrightarrow G^-$, $V^\dagger \rightarrow 1$, and again by reversing the sign in all terms that are odd in the Goldstone fields.

B An alternative flavour basis: real and diagonal Yukawas

There is a certain freedom of choosing the flavour basis for the SMEFT Lagrangian to use for presenting the phenomenological constraints. This freedom stems from the fact that
both the dimension-four Yukawas ($Y_f$) and the dimension-six SMEFT coefficients ($C_{fH}'$) contribute to the observed masses (and mixings) of fermions. As long as this is guaranteed any basis is equivalent. However, since a UV extension of the SM matched to SMEFT would in general induce contributions both to $Y_f$ and $C_{fH}'$ there is no clear notion of a “better” or a more “physical” basis for $Y_f$. We have thus opted to present the bounds in the mass-eigenstate basis as discussed in section 2 (see also ref. [21]). The advantage of this basis is that it is the one required for computations and is also directly related to the κ-framework basis (see discussion around eq. (2.11)). Another approach would be to attempt to provide the constraints in terms of basis-independent, i.e., Jarlskog-type, invariants (see recent ref. [92]).

In this appendix, we discuss a different choice of basis that has been used in the literature [20]. We comment on the differences and point out a consistency condition on the Wilson coefficients in this basis that has been missed in the literature. As before, our starting point is the fermion mass term in eq. (2.3):

$$L_{\text{mass}} = - \sum_{f=u,d,\ell} \frac{v}{\sqrt{2}} \bar{f}_L \left( Y_f - \frac{v^2}{2\Lambda^2} C_{fH}' \right) f_R + \text{h.c.}, \quad (B.1)$$

with $Y_f$ and $C_{fH}'$ generic, complex $3 \times 3$ matrices. Now, instead of diagonalising the full matrix in parentheses, as we did, one may choose to rotate the fermion fields by a biunitary transformation that diagonalises only the dimension-four Yukawa matrices $Y_f$. In other words, we rotate with the transformation matrices

$$f_L \to \tilde{U}_f f_L, \quad f_R \to \tilde{W}_f f_R, \quad (B.2)$$

for $f = u, d, \ell$, such that after the rotation

$$L_{\text{mass}} = - \sum_{u,d,\ell} \frac{v}{\sqrt{2}} \tilde{f}_L \left( \tilde{Y}_f - \frac{v}{2\Lambda^2} \tilde{C}_{fH} \right) f_R + \text{h.c.}, \quad (B.3)$$

where now the matrices $\tilde{Y}_f \equiv \tilde{U}_u Y_u \tilde{W}_f$ are diagonal and real with entries that, however, do not correspond to the observed fermion masses. The kinetic terms of quarks is also affected by the transformation in eq. (B.2). At this stage this means that there is a matrix $\tilde{V} \equiv U_u^\dagger U_d$ in the kinetic terms. Moreover, we have defined $\tilde{C}_{fH} \equiv \tilde{U}_u^\dagger C_{fH} \tilde{W}_f$. In general, the matrices $\tilde{C}_{fH}$ are not diagonal. However, similarly to the discussion of section 2, one can make the UV assumption that they turn out to be complex but diagonal (or ignore off-diagonal entries). This is possible if the matrices $U_f$, $W_f$ simultaneously diagonalise both $Y_f$ and $C_{fH}$. Below we assume that this is the case. The analysis of ref. [20] uses this basis, i.e., $\tilde{C}_{fH \pm}$ or rescalings of them, to present the phenomenological constraints on the SMEFT parameter space.

Since in general, the elements of $\tilde{C}_{fH}$ contain phases, we still have not rotated to the mass-eigenstate basis. To this end we perform an additional (flavour-diagonal) rotation on the right-handed fields as in ref. [20]:

$$f_{R,i} \to e^{i \theta_{f,i}} f_{R,i}, \quad f = u, d, \ell. \quad (B.4)$$
Note that this chiral rotation leaves the kinetic term invariant. The phases $\theta_{f,i}$ are fixed such that they absorb any phase in the corresponding $\hat{C}_{fH}$. Therefore after the rotation

$$L_{\text{mass}} = -\sum_{u,d,\ell} \frac{v}{\sqrt{2}} \tilde{f}_L y^f_{\text{SM}} f_R + \text{h.c.} \quad (B.5)$$

Here, the $y^f_{\text{SM}}$ are diagonal and real matrices with entries that correspond to the observed fermion masses, i.e., we have $m_f = \frac{v}{\sqrt{2}} y^f_{\text{SM}}$, obtained from

$$y^f_{\text{SM}} \equiv \left( \hat{Y}_f - \frac{v^2}{2\Lambda^2} \hat{C}_{fH} \right) e^{i\theta_f} = \hat{U}_f^\dagger \left( \hat{Y}_f - \frac{v^2}{2\Lambda^2} \hat{C}_{fH} \right) \hat{W}_f e^{i\theta_f}. \quad (B.6)$$

The phases $\theta_f$ are thus fixed and can be computed as functions of the physical masses and the $\hat{C}_{fH}$ entries via

$$\sin \theta_{f,i} = \frac{v}{2\sqrt{2}m_{f,i}} \frac{v^2}{\Lambda^2} \hat{C}_{fH-} = \frac{v}{2\sqrt{2}m_{f,i}} \frac{v^2}{\Lambda^2} \text{Im}[\hat{U}_f^\dagger \hat{C}_{fH} \hat{W}_f]. \quad (B.7)$$

As we have the restriction $|\sin \theta_{f,i}| \leq 1$, this gives a consistency bound on $\hat{C}_{fH-}$ or equivalently on $\text{Im}[\hat{U}_f^\dagger \hat{C}_{fH} \hat{W}_f]$.

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