Model-independent Bounds on the Standard Model Effective Theory from Flavour Physics

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Abstract

Meson-antimeson mixing provides the most stringent constraints on baryon- and lepton-number conserving New Physics, probing scales higher than $10^5$ TeV. In the context of the effective theory of weak interactions, these constraints translate into severe bounds on the coefficients of $\Delta F = 2$ operators. Generalizing to the effective theory invariant under the Standard Model gauge group, valid above the electroweak scale, the bounds from $\Delta F = 2$ processes also affect $\Delta F = 1$ and even $\Delta F = 0$ operators, due to log-enhanced radiative corrections induced by Yukawa couplings. We systematically analyze the effect of the renormalization group evolution above the electroweak scale and provide for the first time the full set of constraints on all relevant dimension-six operators.

Keywords:
Flavour Changing Neutral Currents, New Physics, Standard Model Effective Field Theory

The absence of tree-level flavour changing neutral currents (FCNC) in the Standard Model (SM), and their suppression at the loop level [1, 2, 3], make FCNC processes a very sensitive probe of New Physics (NP). In particular, meson-antimeson mixing, a FCNC process with flavour quantum number $F$ changed by two units ($\Delta F = 2$), provides to date the most stringent constraints on baryon- and lepton-number conserving NP, reaching an astonishing NP scale of $\mathcal{O}(10^5)$ TeV for strongly-interacting NP with arbitrary flavour structure [4, 5, 6]. This extraordinary NP sensitivity is due both to the Glashow-Iliopoulos-Maiani (GIM) mechanism and to the hierarchical structure of quark masses and mixing angles. Indeed, the bound on the NP scale can be lowered to a few TeV by requiring Minimal Flavour Violation (MFV), i.e. the absence of new sources of flavour violation beyond Yukawa couplings [7, 8]. In the MFV case, the sensitivity becomes comparable to other indirect probes of NP such as electroweak (EW) precision observables or Higgs signal strengths, see, e.g., the recent works in [9, 10].

If NP arises at scales much higher than the EW one, its leading effects in the EW and flavour sectors can be parameterized in terms of dimension-six local operators built of SM fields and invariant under the SM gauge group. Those operators, together with the SM, form the so-called Standard Model effective field theory (SMEFT) [11, 12]. Quantum corrections due to SM interactions induce a renormalization group (RG) run-
ning of SMEFT operators, which can generate $\Delta F = 2$ operators starting from $\Delta F = 1$ ones. One may then wonder if mixing onto $\Delta F = 2$ operators implies any relevant bound on $\Delta F = 1$ ones, to be eventually compared with present constraints from $\Delta F = 1$ transitions and/or from other EW processes. RG effects also modify the Yukawa couplings, leading to a mismatch between the flavour properties of the SMEFT at the NP and EW scales. This effect is relevant for non-universal operators in the SMEFT, since gauge invariance prevents a full alignment in flavour space of non-universal operators involving left-handed doublets, leading unavoidably to FCNC contributions that depend on the flavour structure at the NP scale. In this Letter we move our first steps towards a deep investigation of the flavour structure of the SMEFT, focusing on $\Delta F = 2$ transitions, the most sensitive probes of NP in the flavour sector, and present bounds on all relevant operators, including all leading RG effects.

Let us begin highlighting the importance of $\Delta F = 2$ processes. Extending the notation of [8], we define the fundamental FCNC MFV coupling between generations $i$ and $j$ in the basis of diagonal down or up Yukawa couplings as $(\lambda^d_{FC})_{i\neq j} \equiv (Y_U^c)_{ij} \sim Y_t^2 V^*_{3j} V_{3i}$ or $(\lambda^u_{FC})_{i\neq j} \equiv (Y_D^c)_{ij} \sim Y_t^2 V^*_{33} V_{3j}$ respectively; $i, j = 1, 2, 3$ are flavour indices, $Y_q$ the Yukawa coupling for quark $q$ in the diagonal basis and $V$ the CKM matrix. We characterize the typical SM FCNC scale as $\Lambda_0 \equiv Y_1 \sin^2 \theta_W M_W / \alpha \sim 2.3$ TeV. An MFV-type NP model will generate $\Delta F = 1$ and $\Delta F = 2$ operators with the same chiral structure as the SM contribution, up to corrections proportional to further powers of Yukawa couplings, with coefficients of $O((\lambda^d_{FC})^2 / \Lambda^2)$ for $\Delta F = 2$ and of $O((\lambda^d_{FC})_{ij}^2 / \Lambda^2)$ for $\Delta F = 1$ processes in the down and up sector respectively. In the down sector, those have the same structure of top-mediated SM contributions, so the most stringent constraints are expected from top-dominated processes such as meson-antimeson mixing or $b \to s \gamma$, leading to lower bounds on the NP scale of few TeV [7, 8]. For a generic model, constraints get much more severe due not only to the absence of the SM CKM and GIM suppression, but also to the possible presence of right-handed flavour changing neutral currents, which are both enhanced by RG evolution and by hadronic matrix elements [14, 15, 16, 17]. One expects to constrain the ratio of the NP coefficients over the NP scale $\Lambda$ as follows:

$$C_{\Delta F=2}^{NP} / \Lambda^2 < \epsilon_{\Delta F=2} C_{\Delta F=2}^{SM} / \Lambda_0^2,$$  

where for a short-distance-dominated meson mixing process $C_{\Delta F=2}^{SM} \sim (\lambda^d_{FC})_{ij}^2$, while $\epsilon_{\Delta F=1}$ is the experimentally allowed fraction of NP contributions to the $\Delta F = 1$ process times the ratio of SM over NP matrix elements. For instance, $\Delta S = 2$ operators with mixed chirality yield $\epsilon_{\Delta F=2} < 10^{-2}$, and plugging in eq. (1) $(\lambda^d_{FC})_{ij}^2 \sim 10^{-4}$ and $C_{\Delta S=2} \sim 1$, a bound of $\Lambda \gtrsim 10^5$ TeV can be obtained from CP violation measurements in the kaon system, see e.g. [18]. We can now estimate the importance of the running from $\Lambda$ to $M_W$ that turns a $\Delta F = 1$ operator into a $\Delta F = 2$ one. We introduce the RG factor $\mathcal{R} \equiv \log(\Lambda / M_W)/(16\pi^2)$, that is order of percent for NP scales above the TeV. Inspired by eq. (1), we naively estimate the lower limit on $\Lambda$ from the mixing into $\Delta F = 2$ to be of order

$$\Lambda^2 \gtrsim C_{\Delta F=1}^{NP} \frac{\mathcal{R} \Lambda_0^2}{\epsilon_{\Delta F=2} (\lambda^d_{FC})_{ij}^2}.$$  

From eq. (2) and the bound one analogously estimates for $\Delta F = 1$ transitions, it follows that $\Delta F = 2$ constraints overcome $\Delta F = 1$ ones if $\mathcal{R} > \epsilon_{\Delta F=2}/\epsilon_{\Delta F=1}$. This is expected to be the case for a large class of operators, since $\epsilon_{\Delta F=2}$ is typically at the percent level, in particular for limits regarding CP violation, while $\epsilon_{\Delta F=1}$ can easily be of $O(1)$ or even (much) larger for two reasons: i) in the SM, the MFV-type top contribution to many $\Delta F = 1$ processes is not dominant with respect to charm or light quarks; ii) the calculation of the relevant matrix elements turns out to be much more uncertain, often plagued by long-distance contributions. Having argued on the general relevance of $\Delta F = 2$ constraints, we present in the following the results of our study
Table 1: Bounds on $X_{1,4,5}^{(R)}$ in GeV$^{-2}$ from the analysis of refs. [5, 13]. See the text for details.

| $ij$ | $X_{1}^{(R)}$ | $X_{1}^{(I)}$ | $X_{4}^{(R)}$ | $X_{4}^{(I)}$ | $X_{5}^{(R)}$ | $X_{5}^{(I)}$ |
|------|----------------|----------------|----------------|----------------|----------------|----------------|
| $s \leftrightarrow d$ | $2.5 \cdot 10^{-13}$ | $2.0 \cdot 10^{-11}$ | $2.0 \cdot 10^{-12}$ | $2.0 \cdot 10^{-11}$ | $2.0 \cdot 10^{-12}$ | $2.0 \cdot 10^{-13}$ |
| $c \leftrightarrow u$ | $5.9 \cdot 10^{-14}$ | $2.0 \cdot 10^{-14}$ | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ |
| $b \leftrightarrow s$ | $9.5 \cdot 10^{-12}$ | $4.0 \cdot 10^{-12}$ | $4.0 \cdot 10^{-12}$ | $4.0 \cdot 10^{-12}$ | $4.0 \cdot 10^{-12}$ | $4.0 \cdot 10^{-12}$ |

Table 2: Constraints on (real, imaginary) parts of Wilson coefficients $C_{ij}^{H(1,3)}$ obtained from $K \rightarrow \bar{K}$ (⊗), $D \rightarrow D$ (△), $B_d \rightarrow B_{d} (\triangle)$ or $B_s \rightarrow B_{s} (\triangledown)$ mixing. Middle and right columns correspond to flavour alignment along the down-quark (up-quark) sector. Entries with no bound are denoted by ⊗. A single bound is quoted for entries required to be real by Hermiticity.

| $ijkl$ | $Y_D$ diag | $Y_U$ diag |
|--------|----------------|----------------|
| 2221   | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ |
| 2222   | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ |
| 2223   | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ |
| 3321   | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ |
| 3322   | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ |
| 3331   | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ |
| 3332   | $2.0 \cdot 10^{-13}$ | $2.0 \cdot 10^{-13}$ |

Table 3: Same as table 2 for $C_{ij}^{LeQu}$ and $C_{ij}^{LedQ}$.

and compare them to a few examples from the literature. We collect our main findings in tables 2-7. A more general study accounting for all other bounds coming from $\Delta F = 1$ measurements and EW precision tests will be given elsewhere. Previous work in this direction can be found in [19, 20, 21, 22, 23, 24, 25].

The most general $\Delta F = 2$ weak effective Hamiltonian (see eqs. (6)-(7) of [5] and also [26, 16]) matches at tree level only five types of gauge-invariant operators [27, 28]:

$$C_1(\mu W) = -\left( C^{QQ(1)}(\mu W) + C^{QQ(3)}(\mu W) \right) / \Lambda^2$$

$$C_1(\mu W) = -C^{qq}(\mu W) / \Lambda^2$$

$$C_4(\mu W) = C^{QQ(5)}(\mu W) / \Lambda^2$$

$$C_5(\mu W) = \left( 2 C^{Q(1)}(\mu W) - \frac{1}{3} C^{QQ(8)}(\mu W) \right) / \Lambda^2$$

For our notation on the SMEFT, see table 8 and also [29, 30]. $Q$ and $q = u, d$ represent

Table 4: Same as table 2 for $C_{ijkl}^{(1)}$ and $C_{ijkl}^{(8)}$. 

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Table 6: Same as table 2 for $C_{ijkl}^{QQ(1;3)}$ and $C_{ijkl}^{dQd}$. 

$SU(2)_L$ gauge doublet and singlets respectively, and $\mu_W \simeq M_W$ is the matching scale. Flavour indices in eq. (3) have been understood. The Wilson coefficients of the SMEFT evolve according to coupled RG equations that we solve at lowest order keeping the logarithmic term:

$$C_A(\mu_W) \simeq (\delta_{AB} - \beta_{AB} \mathcal{R}) C_B(\Lambda),$$  

where $A$ and $B$ are generalized indices running over the flavour structure and field content of the effective theory, and $\beta_{AB}$ characterizes the quantum mixing of $B$ into $A$. Eqs. (3)-(4) allow us to translate existing bounds of the form $(X^{(R)}, X^{(I)})$ on the $\Delta F = 2$ operators of the weak effective theory, see table 1, into constraints on the SMEFT.
barring accidental cancellations:

$$|C_B^{(R)}(\Lambda)| < \left( \frac{\Lambda^2 X_A^{(R)}}{|\delta_{AB} - \beta_{AB}^{(R)}|} \right) \left( \frac{\Lambda^2 X_A^{(I)}}{|\beta_{AB}^{(I)}|} \right),$$

where (R,I) superscripts denote real and imaginary parts. While the bounds obtained through eq. (5), i.e. switching on the real or imaginary

$$|C_B^{(I)}(\Lambda)| < \left( \frac{\Lambda^2 X_A^{(I)}}{|\delta_{AB} - \beta_{AB}^{(R)}|} \right) \left( \frac{\Lambda^2 X_A^{(R)}}{|\beta_{AB}^{(I)}|} \right),$$

Table 7: Same as table 2 for $C^{Q(d)}_{ijkl}$, $C^{Q(d)}_{ijkl}$, $C^{Q(u)}_{ijkl}$ and $C^{Q(u)}_{ijkl}$.

Table 8: Operators involved in the analysis. $H$, $Q$, and $L$ are Higgs, quark, and lepton weak doublets; $d$, $u$ and $e$ quark and lepton $SU(2)_L$ singlets. $T^{\pm}_{\alpha=1...8}$ are $SU(3)_c$ generators; $\tau^{A=1,2,3}$ $SU(2)_L$ ones. $\mu$ and $j,k,l,m$ are Lorentz and flavour indices.
part of a single Wilson coefficient at a time, are generally valid, a model-dependent analysis becomes mandatory whenever two or more Wilson coefficients (including real and imaginary parts) are close to the bounds we provide, since interference effects might become important in that case.

A few remarks are in order. The state-of-the-art knowledge on the anomalous dimensions in eq. (4) encodes all leading order effects in the Higgs self-coupling, Yukawa insertions and SM gauge couplings, see [31, 32, 33, 34]. For our purpose, it suffices to consider the RG effects induced by the Yukawa matrices, including also their (non-negligible) change from the high scale $\Lambda$ to $\mu_W$, and by the strong gauge coupling. Entries for the CKM matrix are taken from the NP fit of [13], neglecting small corrections expected in the SMEFT context [35] which would produce effects of $O(\Lambda^{-4})$ in our analysis. All other SM parameters are defined at the EW scale following [36, 37].

In order to accurately study CP violation effects as well, we developed a numerical code enabling us to fully explore the case of complex-valued Wilson coefficients. The starting point for our numerical analysis is the update [13] of the study in [5], which provides bounds on the coefficients of the weak effective Hamiltonian. The $(X^{(R)}, X^{(I)})$ bounds we present are obtained from the (symmetrized) 95% probability interval per single Wilson coefficient of the $\Delta F = 2$ weak Hamiltonian at $\mu_W$. We run up to the NP scale $\Lambda = 1 \text{ TeV}$ and present bounds on the SMEFT Wilson coefficients at that scale. We discard values larger than $4\pi$ for $\Lambda = 1 \text{ TeV}$. One can convert a bound $C_A < X$ at $1 \text{ TeV}$ into a bound on $\Lambda = 1/\sqrt{X}$ for $C_A = 1$, up to terms of $O(\log(\Lambda/\text{TeV}))$.

In table 2 we collect the bounds on the coefficients of operators involving Higgs and left-handed quark bilinears. Comparing with [38, 39], our bounds on non-flavour-universal coefficients are two to three orders of magnitude stronger than the ones on flavour-universal operators coming from EW precision data and Higgs physics. Remarkably, what reported in table 2 is a pure outcome of the RG flow in the context of the SMEFT. The same is true for table 3, where we present the constraints on the SMEFT dimension-six operator involving anti-symmetric contractions of quark and lepton weak doublets, together with quark and charged lepton singlets. The operator $O^{L_{eQ}u}$ mixes into left-right $\Delta F = 2$ operators in the up sector via the up-quark Yukawa matrix and the diagonal lepton one, $Y_\ell$. Consequently, it gets constrained only by $D - \bar{D}$ mixing in the basis where $Y_D$ is diagonal and $Y_U$ can generate flavour mixing. Conversely, operator $O^{L_{eQ}d}$ mixes into left-right $\Delta F = 2$ operators in the down sector via $Y_t$ and $Y_D$, getting constrained only in the basis of diagonal $Y_U$. In table 4 we present bounds on four-quark operators involving both up- and down-type singlets, which also mix into flavour-violating $\Delta F = 2$ operators in the down and up sectors only via the operator mixing generated by $Y_U$ and $Y_D$ respectively. Another class of operators contributing to $\Delta F = 2$ processes only through mixing via Yukawa interactions is given by four-quark operators involving two doublets, an up-type singlet and a down-type singlet. The set of bounds found is reported in table 5.

Let us turn now to the operators appearing in eq. (3). Depending on the flavour indices, we may handle two different situations: i) the operator matches directly onto the $\Delta F = 2$ operators at the EW scale, after going to the mass eigenstate basis; ii) the operator mixes with the $\Delta F = 2$ ones only via RG flow. Note that for i) one still needs to keep track of RG effects inducing a misalignment of the Yukawa couplings between the cutoff and the EW scale. We stress here how $D - \bar{D}$ and $K - \bar{K}$ mixings play a crucial role in constraining operators involving quark left-handed doublets. CKM misalignment between $Y_U$ and $Y_D$ ensures indeed a non-vanishing contribution of operators $O^{QQ}_{1,3}$ either to $\Delta S = 2$ or $\Delta C = 2$ transitions, leading to stringent constraints from at least one of the two processes. The strongest constraints on $O^{dd}$ come instead from purely right-handed operators and consequently do not depend on the alignment in flavour space once the aforementioned effects from Yukawa running are taken into account. Bounds for operators $O^{QQ}_{1,3,dd}$ are in table 6 and are strictly related to i). For $O^{uu}$ constraints follow mainly from ii) and we find: $C^{uu}_{1112} < (\mathcal{O},6.1\,^\circ)$ TeV$^{-2}$ and
we collect the large set of constraints found for $O^{Q_d(1,8), Q_u(1,8)}$. The discussion of the $\Delta F = 2$ bounds on these operators is analogous to the one already given for $O^{Q_Q(1,3)}$.

The results derived so far represent a very serious challenge for models with new sources of flavour violation, including those cases where NP couples solely (or dominantly) to third generation quarks. Our strong constraints can be considerably weakened in MFV models. Within MFV scenarios, we constrain $O^{H_Q(1,3)}, O^{Q_Q(1,3)}, O^{Q_u(1)}$ and $O^{Q_u(8)}$ operators, probing NP scales as low as 1.3, 9.9, 1.4 and 0.8 TeV respectively. These may be regarded as the weakest possible limits to date on heavy new dynamics coupled to the quark sector.

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