Flavor Physics Constraints  
for Physics Beyond the Standard Model  

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Abstract  

In the last decade, huge progress in experimentally measuring and theoretically understanding flavor physics has been achieved. In particular, the accuracy in the determination of the CKM elements has been greatly improved, and a large number of flavor changing neutral current processes, involving $b \to d$, $b \to s$ and $c \to u$ transitions, and of CP violating asymmetries, have been measured. No evidence for new physics has been established. Consequently, strong constraints on new physics at high scale apply. In particular, the flavor structure of new physics at the TeV scale is strongly constrained. We review these constraints and we discuss the future prospects to better understand the flavor structure of physics beyond the Standard Model.
I. INTRODUCTION

The term “flavors” is used, in the jargon of particle physics, to describe several copies of the same gauge representation, namely several fields that are assigned the same quantum charges. Within the Standard Model (SM), when thinking of its unbroken $SU(3)_C \times U(1)_{EM}$ gauge group, there are four different types of particles, each coming in three flavors:

- Up-type quarks in the $(3)_{+2/3}$ representation: $u, c, t$;
- Down-type quarks in the $(3)_{-1/3}$ representation: $d, s, b$;
- Charged leptons in the $(1)_{-1}$ representation: $e, \mu, \tau$;
- Neutrinos in the $(1)_{0}$ representation: $\nu_1, \nu_2, \nu_3$.

The term “flavor physics” refers to interactions that distinguish between flavors. By definition, gauge interactions, namely interactions that are related to unbroken symmetries and mediated therefore by massless gauge bosons, do not distinguish among the flavors and do not constitute part of flavor physics. Within the Standard Model, flavor-physics refers to the weak and Yukawa interactions. With New Physics (NP), there are likely to be additional ‘flavored’ interactions.

The term “flavor changing” refers to processes where the initial and final flavor-numbers (that is, the number of particles of a certain flavor minus the number of anti-particles of the same flavor) are different. In “flavor changing charged current” processes, both up-type and down-type flavors, and/or both charged lepton and neutrino flavors are involved. Examples are (i) muon decay via $\mu \rightarrow e\bar{\nu}_i \nu_j$, and (ii) $K^- \rightarrow \mu^- \bar{\nu}_j$ (which corresponds, at the quark level, to $s\bar{u} \rightarrow \mu^- \bar{\nu}_j$). Within the Standard Model, these processes are mediated by the $W$-bosons and occur at tree level. In “flavor changing neutral current” (FCNC) processes, either up-type or down-type flavors but not both, and/or either charged lepton or neutrino flavors but not both, are involved. Example are (i) muon decay via $\mu \rightarrow e\gamma$ and (ii) $K_L \rightarrow \mu^+ \mu^-$ (which corresponds, at the quark level, to $s\bar{d} \rightarrow \mu^+ \mu^-$). Within the Standard Model, these processes do not occur at tree level, and are often highly suppressed. This situation makes FCNC particularly sensitive to new physics: If the new physics does not have the same flavor suppression factors as the standard model, then it could contribute to FCNC comparably to the standard model even if it takes places at energy scales that are orders of magnitude higher than the weak scale. The fact that flavor physics is a very sensitive probe of high energy physics is the main reason for the experimental effort to measure flavor parameters and the theoretical effort to interpret these data.
II. THE STANDARD MODEL

The Standard Model (SM) is defined as follows:

(i) The gauge symmetry is

\[ G_{\text{SM}} = SU(3)_C \times SU(2)_L \times U(1)_Y. \]  

It is spontaneously broken by the vacuum expectation values (VEV) of a single Higgs scalar, \( \phi(1, 2) \) \( \frac{1}{2} \) \( (\langle \phi^0 \rangle = v/\sqrt{2}) \):

\[ G_{\text{SM}} \to SU(3)_C \times U(1)_{\text{EM}}. \]  

(ii) There are three fermion generations, each consisting of five representations of \( G_{\text{SM}} \):

\[ Q_{Li}(3, 2)_{1/6}, \ U_{Ri}(3, 1)_{+2/3}, \ D_{Ri}(3, 1)_{-1/3}, \ L_{Li}(1, 2)_{-1/2}, \ E_{Ri}(1, 1)_{-1}. \]  

The Standard Model Lagrangian, \( \mathcal{L}_{\text{SM}} \), is the most general renormalizable Lagrangian that is consistent with the gauge symmetry (2.1), the particle content (2.3) and the pattern of spontaneous symmetry breaking (2.2). It can be divided to three parts:

\[ \mathcal{L}_{\text{SM}} = \mathcal{L}_{\text{kinetic+gauge}} + \mathcal{L}_{\text{Higgs}} + \mathcal{L}_{\text{Yukawa}}. \]  

The source of all flavor physics is in the Yukawa interactions:

\[ -\mathcal{L}_{\text{Yukawa}} = Y^d_{ij} \overline{Q}_{Li} \phi D_{Rj} + Y^u_{ij} \overline{Q}_{Li} \tilde{\phi} U_{Rj} + Y^e_{ij} \overline{L}_{Li} \phi E_{Rj} + \text{h.c.}. \]  

(where \( \tilde{\phi} = i\tau_2 \phi^\dagger \)). This part of the Lagrangian is, in general, flavor-dependent (that is, \( Y^f \propto 1 \)) and CP violating.

In the absence of the Yukawa matrices \( Y^d \), \( Y^u \) and \( Y^e \), the SM has a large \( U(3)^5 \) global symmetry:

\[ G_{\text{global}}(Y^{u,d,e} = 0) = SU(3)_q^3 \times SU(3)_l^2 \times U(1)^5. \]  

The non-Abelian part of this symmetry is particularly relevant to flavor physics:

\[ SU(3)_q^3 = SU(3)_Q \times SU(3)_U \times SU(3)_D, \]
\[ SU(3)_l^2 = SU(3)_L \times SU(3)_E. \]  

The point that is important for our purposes is that \( \mathcal{L}_{\text{kinetic+gauge}} + \mathcal{L}_{\text{Higgs}} \) respect the non-Abelian flavor symmetry \( S(3)_q^3 \times SU(3)_l^2 \), under which

\[ Q_L \to V_Q Q_L, \ U_R \to V_U U_R, \ D_R \to V_D D_R, \ L_L \to V_L L_L, \ E_R \to V_E E_R, \]  

where \( V_Q, V_U, V_D, V_L, V_E \) are 3x3 unitary matrices that parametrize the flavor changes. These matrices must satisfy the unitarity constraint:

\[ V_Q^\dagger V_Q = V_U^\dagger V_U = V_D^\dagger V_D = V_L^\dagger V_L = V_E^\dagger V_E = I. \]  

This condition ensures that the transition matrices are well-defined.
where the $V_i$ are unitary matrices. The Yukawa interactions (2.5) break the global symmetry,

$$G_{\text{global}}(Y_{u,d,e} \neq 0) = U(1)_B \times U(1)_e \times U(1)_\mu \times U(1)_\tau.$$  

(2.9)

(Of course, the gauged $U(1)_Y$ also remains a good symmetry.) Thus, the transformations of Eq. (2.8) are not a symmetry of $\mathcal{L}_{\text{SM}}$. Instead, they correspond to a change of the interaction basis. These observations also offer an alternative way of defining flavor physics: it refers to interactions that break the $SU(3)^5$ symmetry (2.8). Thus, the term “flavor violation” is often used to describe processes or parameters that break the symmetry.

Using the transformation (2.8), one can choose an interaction basis where the number of parameters is minimized. A useful example of such a basis for the quark Yukawa matrices is the following:

$$Y^d = \lambda_d, \quad Y^u = V^\dagger \lambda_u,$$  

(2.10)

where $\lambda_{d,u}$ are diagonal,

$$\lambda_d = \text{diag}(y_d, y_s, y_b), \quad \lambda_u = \text{diag}(y_u, y_c, y_t),$$  

(2.11)

while $V$ is a unitary matrix that depends on three real angles and one complex phase. We conclude that there are 10 quark flavor parameters: 9 real ones and a single phase. In the mass basis, one identifies six of the real parameters as the six quark masses, while the remaining three real and one imaginary parameters appear in the CKM matrix $V$, which describes the couplings of the charged weak-force carriers, the $W^{\pm}$-bosons, with quark-antiquark pairs.

Within the standard model, all flavor changing physics comes from the $V$ matrix. There are various ways to choose the four parameters of $V$. One of the most convenient ways is given by the Wolfenstein parametrization, where the four mixing parameters are $(\lambda, A, \rho, \eta)$ with $\lambda = |V_{us}| = 0.23$ playing the role of an expansion parameter and $\eta$ representing the CP violating phase:

$$V \simeq \left( \begin{array}{ccc} 1 - \frac{1}{2} \lambda^2 - \frac{1}{8} \lambda^4 & \lambda & A \lambda^3 (\rho - i \eta) \\ -\lambda + \frac{1}{2} A^2 \lambda^5 [1 - 2(\rho + i \eta)] & 1 - \frac{1}{2} \lambda^2 - \frac{1}{8} \lambda^4 (1 + 4 A^2) & A \lambda^2 \\ A \lambda^3 [1 - (1 - \frac{1}{2} \lambda^2)(\rho + i \eta)] & -A \lambda^2 + \frac{1}{2} A \lambda^4 [1 - 2(\rho + i \eta)] & 1 - \frac{1}{2} A^2 \lambda^4 \end{array} \right).$$  

(2.12)

One can assume that flavor changing processes are fully described by the SM, and check the consistency of the various measurements with this assumption. The values of $\lambda$ and $A$ are known rather accurately from, respectively, $K \to \pi \ell \nu$ and $b \to c \ell \nu$ decays:

$$\lambda = 0.2257 \pm 0.0010, \quad A = 0.814 \pm 0.022.$$  

(2.13)
Then, one can express all the relevant observables as a function of the two remaining parameters, $\rho$ and $\eta$, and check whether there is a range in the $\rho-\eta$ plane that is consistent with all measurements. The list of observables includes the following:

- The rates of inclusive and exclusive charmless semileptonic $B$ decays depend on $|V_{ub}|^2 \propto \rho^2 + \eta^2$;
- The CP asymmetry in $B \to \psi K_S$, $S_{B \to \psi K} = \sin 2\beta = \frac{2\eta(1-\rho)}{(1-\rho)^2 + \eta^2}$;
- The rates of various $B \to DK$ decays depend on the phase $\gamma$, where $e^{i\gamma} = \rho + i\eta \sqrt{\rho^2 + \eta^2}$;
- The rates of various $B \to \pi\pi, \rho\pi, \rho\rho$ decays depend on the phase $\alpha = \pi - \beta - \gamma$;
- The ratio between the mass splittings in the neutral $B$ and $B_s$ systems is sensitive to $|V_{td}/V_{ts}|^2 = \lambda^2[(1-\rho)^2 + \eta^2]$;
- The CP violation in $K \to \pi\pi$ decays, $\epsilon_K$, depends in a complicated way on $\rho$ and $\eta$.

The resulting constraints are shown in Fig. 1.

The consistency of the various constraints is impressive. In particular, the following ranges for $\rho$ and $\eta$ can account for all the measurements:

$$\rho = +0.135^{+0.031}_{-0.016}, \quad \eta = +0.349 \pm 0.017.$$  \hspace{1cm} (2.14)

One can make then the following statement:

Very likely, flavor violation and CP violation in flavor changing processes are dominated by the CKM mechanism.

One can actually go a step further, and allow for arbitrary new physics in all flavor changing processes except for those that have contributions from SM tree diagrams. Then, one can quantitatively constrain the size of new physics contributions to processes such as neutral meson mixing. We do so in the next section.

III. MODEL INDEPENDENT CONSTRAINTS

In order to describe NP effects in flavor physics we can follow two main strategies: (i) build an explicit ultraviolet completion of the model, and specify which are the new fields beyond the SM ones, or (ii) analyse the NP effects using a generic effective-theory approach, or integrating-out the new heavy fields. The first approach is more predictive, but also more model dependent. We
follow this approach in Sect. \( \Box \) and \( \Box \) in two well-motivated SM extensions. In this and the next section we follow the second strategy, which is less predictive but also more general.

Assuming the new degrees to be heavier than SM fields, we can integrate them out and describe NP effects by means of a generalization of the Fermi Theory. The SM Lagrangian becomes the renormalizable part of a more general local Lagrangian which includes an infinite tower of operators with dimension \( d > 4 \), constructed in terms of SM fields, suppressed by inverse powers of an effective scale \( \Lambda > M_W \):

\[
\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \sum c^{(d)}_i \frac{\Lambda^{(d-4)}}{\Lambda} O^{(d)}_i \text{(SM fields)}. \quad (3.1)
\]

This general bottom-up approach allows us to analyse all realistic extensions of the SM in terms of a limited number of parameters (the coefficients of the higher-dimensional operators). The drawback of this method is the impossibility to establish correlations of NP effects at low and high energies: the scale \( \Lambda \) defines the cut-off of the effective theory. However, correlations among different low-
energy processes can still be established implementing specific symmetry properties, such as the
MFV hypothesis (Sect. IV). The experimental tests of such correlations allow us to test/establish
general features of the new theory which holds independently of the dynamical details of the model.
In particular, B, D and K decays are extremely useful in determining the flavor-symmetry breaking
pattern of the NP model.

A. Bounds from $\Delta F = 2$ down-type transitions

The starting points for this analysis is the observation that in several realistic NP models we
can neglect non-standard effects in all cases where the corresponding effective operator is gen-
erated at the tree-level within the SM. This general assumption implies that the experimental
determination of the CKM matrix via tree-level processes is free from the contamination of NP
contributions. Using this determination we can unambiguously predict meson-antimeson mixing
and FCNC amplitudes within the SM and compare it with data, constraining the couplings of the
$\Delta F = 2$ operators in (3.1). Each $\Delta F = 2$ amplitude is then conveniently parametrized in terms of
the shift induced in the modulo and the CPV phase, or the real and imaginary part [10, 11]:

$$\frac{\langle B_q | L_{\text{eff}} | \overline{B}_q \rangle}{\langle B_q | L_{\text{SM}} | \overline{B}_q \rangle} = C_{B_q} e^{2i\phi_{B_q}}, \quad \frac{\text{Re}[\langle K^0 | L_{\text{eff}} | \overline{K}^0 \rangle]}{\text{Re}[\langle K^0 | L_{\text{SM}} | \overline{K}^0 \rangle]} = C_{\Delta m_K} \xrightarrow{\text{Re} \rightarrow \text{Im}} C_{\epsilon_K},$$

(3.2)

An updated analysis of these constraints has been presented in [12] (see Fig. 2). The main conclu-
sions that can be drawn from this analysis can be summarized as follows:

(i) In all the three accessible short-distance amplitudes ($K^0-\overline{K}^0$, $B_d-\overline{B}_d$, and $B_s-\overline{B}_s$) the mag-
nitude of the new-physics amplitude cannot exceed, in size, the SM short-distance contribution.
The latter is suppressed by both the GIM mechanism and the hierarchical structure of the CKM
matrix,

$$A_{\Delta F=2}^{\Delta F=2} \approx \frac{G_F m_t^2}{16\pi^2} (V_{ti}^* V_{lj})^2 \left\langle \overline{M} | (\nabla_{Li} \gamma^\mu Q_{Lj})^2 | M \right\rangle \times F \left( \frac{M_W^2}{m_t^2} \right),$$

(3.3)

where $F$ is a loop function of $O(1)$. As a result, new-physics models with TeV-scale flavored degrees
of freedom and $O(1)$ effective flavor-mixing couplings are ruled out. To set explicit bounds, let us
consider for instance the subset of left-handed $\Delta F = 2$ operators in the generic effective Lagrangian
in (3.1), namely

$$\Delta L^{\Delta F=2} = \sum_{i \neq j} \frac{c_{ij}}{\Lambda^2} (\nabla_{Li} \gamma^\mu Q_{Lj})^2,$$

(3.4)
where the $c_{ij}$ are dimensionless couplings. The condition $|A_{NP}^{F=2}| < |A_{SM}^{F=2}|$ implies

$$\Lambda > \frac{4.4 \text{ TeV}}{|V_{ti}^* V_{tj}|/|c_{ij}|^{1/2}} \sim \begin{cases} 1.3 \times 10^4 \text{ TeV} \times |c_{sd}|^{1/2} \\ 5.1 \times 10^2 \text{ TeV} \times |c_{bd}|^{1/2} \\ 1.1 \times 10^2 \text{ TeV} \times |c_{bs}|^{1/2} \end{cases}$$

(3.5)

The strong bounds on $\Lambda$ for generic $c_{ij}$ of order 1 is a manifestation of what in many specific frameworks (supersymmetry, technicolor, etc.) goes under the name of flavor problem: if we insist that the new physics emerges in the TeV region, we have to conclude that it possesses a highly non-generic flavor structure.

(ii) In the case of $B_d \rightarrow \overline{B}_d$ and $K^0 - \overline{K}^0$ mixing, where both CP conserving and CP-violating observables are measured with excellent accuracy, there is still room for a sizable NP contribution (relative to the SM one), provided that it is to a good extent aligned in phase with the SM amplitude [$O(0.01)$ for the $K$ system and $O(0.3)$ for the $B_d$ system]. This is because the theoretical errors in the observables used to constraint the phases, $S_{B_d \rightarrow \psi K}$ and $\epsilon_K$, are smaller with respect to the theoretical uncertainties in $\Delta m_{B_d}$ and $\Delta m_K$, which constrain the magnitude of the mixing amplitudes.

(iii) In the case of $B_s \rightarrow \overline{B}_s$ mixing, the precise determination of $\Delta m_{B_s}$ does not allow large deviations in modulo with respect to the SM. The constraint is particularly severe if we consider the ratio $\Delta m_{B_s}/\Delta m_{B_d}$, where hadronic uncertainties cancel to a large extent. However, the constraint on the CP-violating phase is quite poor. Present data from CDF [13] and D0 [14] indicate a large
TABLE I: Bounds on representative dimension-six $\Delta F = 2$ operators. Bounds on $\Lambda$ are quoted assuming an effective coupling $1/\Lambda^2$, or, alternatively, the bounds on the respective $c_{ij}$’s assuming $\Lambda = 1$ TeV. Observables related to CPV are separated from the CP conserving ones with semicolons. In the $B_s$ system we only quote a bound on the modulo of the NP amplitude derived from $\Delta m_{B_s}$ (see text). For the definition of the CPV observables in the $D$ system see Ref. [13].

| Operator | Bounds on $\Lambda$ in TeV ($c_{ij} = 1$) | Bounds on $c_{ij}$ ($\Lambda = 1$ TeV) | Observables |
|----------|--------------------------------------|-------------------------------------|-------------|
|          | Re        | Im        | Re        | Im        | $\Delta m_K; \epsilon_K$ |
| $(s_L^\gamma \mu d_L)^2$ | $9.8 \times 10^2$ | $1.6 \times 10^4$ | $9.0 \times 10^{-7}$ | $3.4 \times 10^{-9}$ |
| $(s_R d_L)(\bar s_L d_R)$ | $1.8 \times 10^4$ | $3.2 \times 10^5$ | $6.9 \times 10^{-9}$ | $2.6 \times 10^{-11}$ |
| $(\bar c_L\gamma \mu u_L)^2$ | $1.2 \times 10^3$ | $2.9 \times 10^3$ | $5.6 \times 10^{-7}$ | $1.0 \times 10^{-7}$ |
| $(\bar c_R u_L)(\bar c_L u_R)$ | $6.2 \times 10^3$ | $1.5 \times 10^4$ | $5.7 \times 10^{-8}$ | $1.1 \times 10^{-8}$ |
| $(\bar b_L\gamma \mu d_L)^2$ | $5.1 \times 10^2$ | $9.3 \times 10^2$ | $3.3 \times 10^{-6}$ | $1.0 \times 10^{-6}$ |
| $(\bar b_R d_L)(\bar b_L d_R)$ | $1.9 \times 10^3$ | $3.6 \times 10^3$ | $5.6 \times 10^{-7}$ | $1.7 \times 10^{-7}$ |
| $(\bar b_L\gamma \mu s_L)^2$ | $1.1 \times 10^2$ | | $7.6 \times 10^{-5}$ | $\Delta m_{B_s}$ |
| $(\bar b_R s_L)(\bar b_L s_R)$ | $3.7 \times 10^2$ | | $1.3 \times 10^{-5}$ | $\Delta m_{B_s}$ |

The central value for the CP-violating phase, contrary to the SM expectation. The errors are, however, still large and the disagreement with the SM is at about the 2$\sigma$ level. If the disagreement persists, becoming statistically significant, this would not only signal the presence of physics beyond the SM, but would also rule out a whole subclass of MFV models (see Sect. [V]).

(iv) In $D - \bar D$ mixing we cannot estimate the SM contribution from first principles; however, to a good accuracy this is CP conserving. As a result, strong bounds on possible non-standard CP-violating contributions can still be set. The resulting constraints are only second to those from $\epsilon_K$, and unlike in the case of $\epsilon_K$ are controlled by experimental statistics and could possibly be significantly improved in the near future.

A more detailed list of the bounds derived from $\Delta F = 2$ observables is shown in Table I, where we quote the bounds for two representative sets of dimension-six operators: the left-left operators (present also in the SM) and operators with a different chirality, which arise in specific SM extensions. The bounds on the latter are stronger, especially in the kaon case, because of the larger hadronic matrix elements. The constraints related to CPV correspond to maximal phases, and are subject to the requirement that the NP contributions are smaller than 30% (60%) of the total contributions in the $B_d$ ($K$) system. Since the experimental status of CP violation in the $B_s$ system is not yet settled we simply require that the new physics contributions are smaller than
the observed value of $\Delta m_{B_s}$ (for less naive treatments see e.g. [3, 12]).

B. Correlations between $K$ and $D$ mixing

There are two different features that can provide flavor-related suppression factors: degeneracy and alignment. In general, low energy measurements can only constrain the product of these two suppression factors. An interesting exception occurs, however, for the LL operators of the type \[[3,4]\] where there is an independent constraint on the level of degeneracy \[[10]\]. We here briefly explain this point.

Consider operators of the form

$$\frac{1}{A_{NP}^2} (\overline{Q}_L)(X_Q)_{ij}\gamma_\mu Q_{Lj})(\overline{Q}_L)(X_Q)_{ij}\gamma_\mu Q_{Lj}),$$

where $X_Q$ is an hermitian matrix. Without loss of generality, we can choose to work in the basis defined in Eq. (2.10):

$$Y^d = \lambda_d, \quad Y^u = V^\dagger \lambda_u, \quad X_Q = V_d^\dagger \lambda Q V_d,$$

where $\lambda_Q$ is a diagonal real matrix, and $V_d$ is a unitary matrix which parametrizes the misalignment of the operator (3.6) with the down mass basis.

The experimental constraints that are most relevant to our study come from $K^0-K^0$ and $D^0-D^0$ mixing, which involve only the first two generation quarks. When studying new physics effects, ignoring the third generation is often a good approximation to the physics at hand. Indeed, even when the third generation does play a role, our two generation analysis is applicable as long as there are no strong cancellations with contributions related to the third generation. In a two generation framework, $V$ depends on a single mixing angle (the Cabibbo angle $\theta_c$), while $V_d$ depends on a single angle and a single phase. To understand various aspects of our analysis, it is useful, however, to provisionally set the phase to zero, and study only CP conserving (CPC) observables. We thus have

$$\lambda_Q = \text{diag}(\lambda_1, \lambda_2), \quad V = \begin{pmatrix} \cos \theta_c & \sin \theta_c \\ -\sin \theta_c & \cos \theta_c \end{pmatrix}, \quad V_d = \begin{pmatrix} \cos \theta_d & \sin \theta_d \\ -\sin \theta_d & \cos \theta_d \end{pmatrix}.$$  

It is convenient to define

$$\lambda_{12} = \frac{1}{2}(\lambda_1 + \lambda_2), \quad \delta_{12} = \frac{\lambda_1 - \lambda_2}{\lambda_1 + \lambda_2}, \quad \lambda_{12} = \delta_{12} \lambda_{12}. $$

Thus $\lambda_{12}$ parametrizes the overall, flavor-diagonal suppression of $X_Q$ (in particular, loop factors), $\delta_{12}$ parametrizes suppression that is coming from approximate degeneracy between the eigenvalues.
of $X_Q$, and $\theta_d$ and $\theta_c - \theta_d$ parametrize the suppression that comes from alignment with, respectively, the down and the up sector.

The main point is the following: Alignment can entirely suppress the contribution to either $K^0 - \bar{K}^0$ mixing ($\theta_d = 0$) or $D^0 - \bar{D}^0$ mixing ($\theta_d = \theta_c$) but not to both. Thus, the flavor measurements give a constraint on $\Lambda_{12}$ which reads

$$\Lambda_{12} \leq 3.8 \times 10^{-3} \left( \frac{\Lambda_{NP}}{1 \text{ TeV}} \right).$$

(3.10)

If we switch on the CP violating phase $\gamma$ in $V_d$ then, for $0.03 \lesssim |\sin\gamma| \lesssim 0.98$, we find

$$\Lambda_{12} \leq 4.8 \times 10^{-4} \sqrt{\sin 2\gamma} \left( \frac{\Lambda_{NP}}{1 \text{ TeV}} \right).$$

(3.11)

We learn that, with a loop suppression of order $\lambda_{12} \sim \alpha_2$, the degeneracy should be stronger than 0.02.

C. Top Physics

While the present direct determination of top-quark decay rates is at the $\mathcal{O}(1\%)$ level of accuracy [17], orders of magnitude improvement is expected at the LHC. With 100 fb$^{-1}$ of data, the LHC will be sensitive (at 95% CL) to branching ratios of $\mathcal{O}\left(10^{-5}\right)$ in the $t \to u^i Z/\gamma$ channels [18], where $u^i = u, c$. Within the SM, these channels have branching ratios in the $10^{-13}$ range, so any experimental observation would be a clear sign of physics beyond the SM. The LHC sensitivity to additional $\Delta t = 1$ processes, such as $t \to c G, \phi$, is more limited (see e.g. Ref. [19] and references therein).

The impact of the projected LHC bounds on top flavor violation can be described in a model independent manner. The leading contributions to the above processes are mediated via a set of dimension-six operators [20], which can be classified according to the SM global flavor symmetries [21]. Focusing on the chirality of the quark fields, these can be written as $\mathcal{O}_{A'B'}$, where $A^i, B = L, R$ and $A^i$ denotes top-quark chirality. Bounds from $B$ physics can give severe constraints on the operators with $B = L$ [21]. The situation changes, however, in case of a flavor alignment of these effective operators, such that they are diagonal in the down-type mass basis. In this limit no bounds can be derived on the $\mathcal{O}_{LR}$ and $\mathcal{O}_{RL}$ operators [22]. On the other hand, the $\mathcal{O}_{LL}$ operators cannot be aligned simultaneously with the down- and up-quark mass bases. As a result, combining bounds on $t \to u^i$ with bounds from $b \to s l^+ l^-$ we can get stringent constraints on the operator $\mathcal{O}^{h}_{LL} = \bar{Q}_i \gamma^\mu (X_Q)_{ij} Q_j \left( \phi^i \hat{D}_\mu \phi \right) / \Lambda_t^2$. Normalising the $SU(3)_Q$ adjoint spurion, $X_Q = a_i T_i$, such
that the projection onto the $SU(3)$ generators has a unit norm ($\sum_i a_i^2 = 1$), we can indentify a unique combination which minimizes the combined bound from top and $B$ constraints. Taking into account the projected LHC sensitivity on $t \to u^i$ decays, we can expect to reach bounds of $O(600 \text{ GeV})$ on the scale $\Lambda_t$ \cite{22}, in the absence of a NP signal.

IV. MINIMAL FLAVOR VIOLATION

A very reasonable, although quite pessimistic, set up which avoids the new physics flavor problem is the so-called Minimal Flavor Violation (MFV) hypothesis. Under this assumption, flavor-violating interactions are linked to the known structure of Yukawa couplings also beyond the SM. As a result, non-standard contributions in FCNC transitions turn out to be suppressed to a level consistent with experiments even for $\Lambda \sim$ few TeV. One of the most interesting aspects of the MFV hypothesis is that it can naturally be implemented within the generic effective Lagrangian in Eq. (3.1). Furthermore, SM extensions where the flavor hierarchy is generated at a scale much higher than other dynamical scales tend to flow to the MFV class of models in the infra-red.

The MFV hypothesis consists of two ingredients \cite{23}: (i) a flavor symmetry and (ii) a set of symmetry-breaking terms. The symmetry is nothing but the large global symmetry of the SM Lagrangian in absence of Yukawa couplings shown in Eq. (2.7). Since this global symmetry, and particularly the $SU(3)$ subgroups controlling quark flavor-changing transitions, is already broken within the SM, we cannot promote it to be an exact symmetry of the NP model. Some breaking would appear at the quantum level because of the SM Yukawa interactions. The most restrictive assumption we can make to protect in a consistent way quark-flavor mixing beyond the SM is to assume that $Y^d$ and $Y^u$ are the only sources of flavor symmetry breaking also in the NP model. To implement and interpret this hypothesis in a consistent way, we can assume that $SU(3)^3_q$ is a good symmetry and promote $Y^u,d$ to be non-dynamical fields (spurions) with non-trivial transformation properties under $SU(3)^3_q$:

$$Y^u \sim (3, \bar{3}, 1) , \quad Y^d \sim (3, 1, \bar{3}) .$$

If the breaking of the symmetry occurs at very high energy scales, at low-energies we would only be sensitive to the background values of the $Y$, \textit{i.e.} to the ordinary SM Yukawa couplings. Employing the effective-theory language, an effective theory satisfies the criterion of Minimal Flavor Violation in the quark sector\footnote{The notion of MFV can be extended also to the lepton sector. However, in this case there is not a unique way} if all higher-dimensional operators, constructed from SM and $Y$ fields, are
invariant (formally) under the flavor group $SU(3)^3_q$ \cite{23}. Invariance under CP may or may not be imposed in addition.

According to this criterion one should in principle consider operators with arbitrary powers of the (dimensionless) Yukawa fields. However, a strong simplification arises by the observation that all the eigenvalues of the Yukawa matrices are small, but for the top one (and possibly the bottom one, see later), and that the off-diagonal elements of the CKM matrix are very suppressed. Working in the basis in Eq. (2.10), and neglecting the ratio of light quark masses over the top mass, we have

$$\left[ Y_u(Y_u^u)^\dagger \right]_{i \neq j}^n \approx y_t^{2n}V_{ti}^*V_{tj}. \quad (4.2)$$

As a consequence, including high powers of the the Yukawa matrices amounts only to a redefinition of the overall factor in (1.2) and the the leading $\Delta F = 2$ and $\Delta F = 1$ FCNC amplitudes get exactly the same CKM suppression as in the SM:

$$A(d^i \rightarrow d^j)_{\text{MFV}} = (V_{ti}^*V_{tj})A_{\text{SM}}^{(\Delta F = 1)} \left[ 1 + a_1 \frac{16\pi^2 M_W^2}{\Lambda^2} \right], \quad (4.3)$$

$$A(M_{ij} - \overline{M}_{ij})_{\text{MFV}} = (V_{ti}^*V_{tj})^2A_{\text{SM}}^{(\Delta F = 2)} \left[ 1 + a_2 \frac{16\pi^2 M_W^2}{\Lambda^2} \right], \quad (4.4)$$

where the $A_{\text{SM}}^{(i)}$ are the SM loop amplitudes and the $a_i$ are $O(1)$ real parameters. The $a_i$ depend on the specific operator considered but are flavor independent. This implies the same relative correction in $s \rightarrow d$, $b \rightarrow d$, and $b \rightarrow s$ transitions of the same type.

Within the MFV framework, several of the constraints used to determine the CKM matrix (and in particular the unitarity triangle) are not affected by NP \cite{26}. In this framework, NP effects are negligible not only in tree-level processes but also in a few clean observables sensitive to loop effects, such as the time-dependent CPV asymmetry in $B_d \rightarrow \psi K_{L,S}$. Indeed the structure of the basic flavor-changing coupling in Eq. (4.4) implies that the weak CPV phase of $B_d - \overline{B_d}$ mixing is $\arg[(V_{td}V_{tb}^*)^2]$, exactly as in the SM. This construction provides a natural (a posteriori) justification of why no NP effects have been observed in the quark sector: by construction, most of the clean observables measured at $B$ factories are insensitive to NP effects in the MFV framework.

In Table II we report a few representative examples of the bounds on the higher-dimensional operators in the MFV framework. For simplicity, only leading spurion dependence is shown on the left-handed column. The built-in CKM suppression leads to bounds on the effective scale

\[\text{to define the minimal sources of flavour symmetry breaking if we want to keep track of non-vanishing neutrino masses} \, \cite{24,23}.\]
of new physics not far from the TeV region. These bounds are very similar to the bounds on flavor-conserving operators derived by precision electroweak tests. This observation reinforces the conclusion that a deeper study of rare decays is definitely needed in order to clarify the flavor problem: the experimental precision on the clean FCNC observables required to obtain bounds more stringent than those derived from precision electroweak tests (and possibly discover new physics) is typically in the 1% − 10% range.

Although MFV seems to be a natural solution to the flavor problem, it should be stressed that (i) this is not a theory of flavor (there is no explanation for the observed hierarchical structure of the Yukawas), and (ii) we are still far from having proved the validity of this hypothesis from data (in the effective theory language we can say that there is still room for sizable new sources of flavor symmetry breaking beside the SM Yukawa couplings [28]). A proof of the MFV hypothesis can be achieved only with a positive evidence of physics beyond the SM exhibiting the flavor-universality pattern (same relative correction in $s \rightarrow d$, $b \rightarrow d$, and $b \rightarrow s$ transitions of the same type) predicted by the MFV assumption. While this goal is quite difficult to be achieved, the MFV framework is quite predictive and thus could easily be falsified: in Table III we list some clean MFV predictions which could be falsified by future experiments. Violations of these bounds would not only imply physics beyond the SM, but also a clear signal of new sources of flavor symmetry breaking beyond the Yukawa couplings.

The idea that the CKM matrix rules the strength of FCNC transitions also beyond the SM has become a very popular concept in recent literature and has been implemented and discussed in several works. It is worth stressing that the CKM matrix represents only one part of the problem: a key role in determining the structure of FCNCs is also played by quark masses, or by

| Operator                                                                 | Bound on $\Lambda$ | Observables                          |
|--------------------------------------------------------------------------|---------------------|--------------------------------------|
| $H_D^\dagger (\bar{D}_R Y_d Y_u^\dagger Y_u^\dagger \sigma_{\mu\nu} Q_L) (e F_{\mu\nu})$ | 6.1 TeV             | $B \rightarrow X_s \gamma, B \rightarrow X_s \ell^+ \ell^-$ |
| $\frac{1}{2} (\bar{Q}_L Y_u Y_u^\dagger \gamma_{\mu} Q_L)^2$               | 5.9 TeV             | $\epsilon_K, \Delta m_B, \Delta m_{Bs}$ |
| $H_D^\dagger (\bar{D}_R Y_d Y_u^\dagger Y_u^\dagger \sigma_{\mu\nu} T^u Q_L) (g_s G^u_{\mu\nu})$ | 3.4 TeV             | $B \rightarrow X_s \gamma, B \rightarrow X_s \ell^+ \ell^-$ |
| $(\bar{Q}_L Y_u Y_u^\dagger \gamma_{\mu} Q_L) (\bar{D}_R \gamma_{\mu} F_R)$ | 2.7 TeV             | $B \rightarrow X_s \ell^+ \ell^-, B_s \rightarrow \mu^+ \mu^-$ |
| $i (\bar{Q}_L Y_u Y_u^\dagger \gamma_{\mu} Q_L) (H_U^\dagger D_{\mu} H_U)$  | 2.3 TeV             | $B \rightarrow X_s \ell^+ \ell^-, B_s \rightarrow \mu^+ \mu^-$ |
| $(\bar{Q}_L Y_u Y_u^\dagger \gamma_{\mu} Q_L) (L_L \mu L_L)$              | 1.7 TeV             | $B \rightarrow X_s \ell^+ \ell^-, B_s \rightarrow \mu^+ \mu^-$ |
| $(\bar{Q}_L Y_u Y_u^\dagger \gamma_{\mu} Q_L) (e D_{\mu} F_{\mu\nu})$     | 1.5 TeV             | $B \rightarrow X_s \ell^+ \ell^-$ |

TABLE II: Bounds on the scale of new physics (at 95% C.L.) for some representative $\Delta F = 1$ [27] and $\Delta F = 2$ [12] MFV operators (assuming effective coupling $\pm 1/\Lambda^2$), and corresponding observables used to set the bounds.
| Observable                              | Experiment            | MFV prediction | SM prediction |
|----------------------------------------|-----------------------|----------------|---------------|
| $\beta_s$ from $A_{CP}(B_s \to \psi \phi)$ | $[0.10, 1.44]$ @ 95% CL | 0.04(5)$^*$    | 0.04(2)       |
| $A_{CP}(B \to X_s \gamma)$            | $< 6\%$ @ 95% CL     | $< 0.02^*$     | < 0.01        |
| $B(\bar{B}_d \to \mu^+ \mu^-)$       | $< 1.8 \times 10^{-8}$| $< 1.2 \times 10^{-9}$ | $1.3(3) \times 10^{-10}$ |
| $B(B \to X_s \tau^+ \tau^-)$          | –                     | $< 5 \times 10^{-7}$ | $1.6(5) \times 10^{-7}$ |
| $B(K_L \to \pi^0 \nu \bar{\nu})$     | $< 2.6 \times 10^{-8}$ @ 90% CL | $< 2.9 \times 10^{-10}$ | $2.9(5) \times 10^{-11}$ |

TABLE III: Some predictions derived in the MFV framework. Stars implies that the prediction is not valid in the GMFV case at large tan $\beta$ (see Section IV B).

In this respect, the MFV criterion provides the maximal protection of FCNCs (or the minimal violation of flavor symmetry), since the full structure of Yukawa matrices is preserved. At the same time, this criterion is based on a renormalization-group-invariant symmetry argument, which can be implemented independently of any specific hypothesis about the dynamics of the new-physics framework. This model-independent structure does not hold in most of the alternative definitions of MFV models that can be found in the literature. For instance, the definition of Ref. [29] (denoted constrained MFV, or CMFV) contains the additional requirement that the effective FCNC operators playing a significant role within the SM are the only relevant ones also beyond the SM. This condition is realized only in weakly coupled theories at the TeV scale with only one light Higgs doublet, such as the MSSM with small tan $\beta$ where no large logs, or sizable anomalous dimension are present [30] (see also [31]). It does not hold in several other frameworks, such as Higgsless models, 5D MFV models, or the MSSM with large tan $\beta$.

In NP models where sizable anomalous dimensions are present, the expansion in powers of the Yukawa spurions cannot be truncated at the first non-trivial order. In this limit, denoted as General MFV (GMFV), higher-order terms in the third-generation Yukawa couplings need to be re-summed. As shown in Eq. (4.2), if only the top-quark Yukawa coupling is of order one, this resummation has negligible consequences for flavour-violating processes. (To measure the effect requires that the accuracy on rare $K$ ($D$) decays would become strong enough to probe the contributions related to the charm (strange) quark Yukawa coupling [30]). However, significant differences from a linear expansion may arise if both top- and bottom-quark Yukawa couplings are order one as it happens, for instance, in the large tan $\beta$ regime.
A. MFV at large $\tan \beta$.

If the Yukawa Lagrangian contains more than a single Higgs field, we can still assume that the Yukawa couplings are the only irreducible breaking sources of $SU(3)_q^3$, but we can change their overall normalization.

A particularly interesting scenario is the two-Higgs-doublet model where the two Higgs fields, $\phi_U$ and $\phi_D$, are coupled separately to up- and down-type quarks:

$$-\mathcal{L}_{\text{Yukawa}}^{2\text{HDM}} = Y^d_{ij} \bar{Q}_{Li} \phi_D D R_j + Y^u_{ij} \bar{Q}_{Li} \phi_U U R_j + Y^e_{ij} \bar{L}_{Li} \phi_D E R_j + \text{h.c.}$$  \hspace{1cm} (4.5)

This Lagrangian is invariant under an extra $U(1)$ symmetry with respect to the one-Higgs Lagrangian in Eq. (2.5): a symmetry under which the only charged fields are $D_R$ and $E_R$ (charge +1) and $\phi_D$ (charge $-1$). This symmetry, denoted $U(1)_{\text{PQ}}$, prevents tree-level FCNCs and implies that $Y^{u,d}$ are the only sources of $SU(3)_q^3$ breaking appearing in the Yukawa interaction (similar to the one-Higgs-doublet scenario). Consistently with the MFV hypothesis, we can then assume that $Y^{u,d}$ are the only relevant sources of $SU(3)_q^3$ breaking appearing in all the low-energy effective operators. This is sufficient to ensure that flavor-mixing is still governed by the CKM matrix, and naturally guarantees a good agreement with present data in the $\Delta F = 2$ sector. However, the extra symmetry of the Yukawa interaction allows us to change the overall normalization of $Y^{u,d}$ with interesting phenomenological consequences in specific rare modes. These effects are related only to the large value of bottom Yukawa, and indeed can be found also in other NP frameworks where there is no extended Higgs sector, but the bottom Yukawa coupling is of order one [30].

Assuming the Lagrangian in Eq. (4.5), the normalization of the Yukawa couplings is controlled by the ratio of the vacuum expectation values of the two Higgs fields, or by the parameter

$$\tan \beta = \langle \phi_U \rangle / \langle \phi_D \rangle .$$  \hspace{1cm} (4.6)

For $\tan \beta \gg 1$ the smallness of the $b$ quark and $\tau$ lepton masses can be attributed to the smallness of $1 / \tan \beta$ rather than to the corresponding Yukawa couplings. As a result, for $\tan \beta \gg 1$ we cannot anymore neglect the down-type Yukawa coupling. Moreover, the $U(1)_{\text{PQ}}$ symmetry cannot be exact: it has to be broken at least in the scalar potential in order to avoid the presence of a massless pseudoscalar Higgs. Even if the breaking of $U(1)_{\text{PQ}}$ and $SU(3)_q^3$ are decoupled, the presence of $U(1)_{\text{PQ}}$ breaking sources can have important implications on the structure of the Yukawa interaction, especially if $\tan \beta$ is large [23, 32–34]. We can indeed consider new dimension-four operators such as

$$\epsilon \bar{Q}_L Y^d D R \tilde{\phi}_U \quad \text{or} \quad \epsilon \bar{Q}_L Y^u Y^u Y^d D R \tilde{\phi}_U ,$$  \hspace{1cm} (4.7)
where $\epsilon$ denotes a generic MFV-invariant $U(1)_{PQ}$-breaking source. Even if $\epsilon \ll 1$, the product $\epsilon \times \tan \beta$ can be $O(1)$, inducing large corrections to the down-type Yukawa sector:

$$
\epsilon \frac{Q_L}{\tan \beta} \rightarrow \epsilon \frac{Q_L}{\tan \beta} \langle \phi_U \rangle = (\epsilon \times \tan \beta) \frac{Q_L}{\tan \beta} \langle \phi_D \rangle.
$$

(4.8)

Since the $b$-quark Yukawa coupling becomes $O(1)$, the large-$\tan \beta$ regime is particularly interesting for helicity-suppressed observables in $B$ physics. One of the clearest phenomenological consequences is a suppression (typically in the $10 - 50\%$ range) of the $B \rightarrow \ell \nu$ decay rate with respect to its SM expectation [35]. Potentially measurable effects in the $10 - 30\%$ range are expected also in $B \rightarrow X_s \gamma$ [36] and $\Delta M_{B_s}$ [37]. Given the present measurements of $B \rightarrow \ell \nu$, $B \rightarrow X_s \gamma$, and $\Delta M_{B_s}$, none of these effects seems to be favored by data. However, present errors are still sizable compared to the estimated NP effects.

The most striking signature could arise from the rare decays $B_{s,d} \rightarrow \ell^+ \ell^-$, whose rates could be enhanced over the SM expectations by more than one order of magnitude [38]. An enhancement of both $B_s \rightarrow \ell^+ \ell^-$ and $B_d \rightarrow \ell^+ \ell^-$ respecting the MFV relation $\Gamma(B_s \rightarrow \ell^+ \ell^-) / \Gamma(B_d \rightarrow \ell^+ \ell^-) \approx |V_{ts} / V_{td}|^2$ would be an unambiguous signature of MFV at large $\tan \beta$ [27].

**B. MFV with additional flavor-diagonal phases**

The breaking of the $SU(3)_q^3$ flavor group and the breaking of the discrete CP symmetry are not necessarily related, and we can add flavor-diagonal CPV phases to generic MFV models [39–41]. Because of the experimental constraints on electric dipole moments (EDMs), which are generally sensitive to such flavour-diagonal phases [40], in this more general case the bounds on the scale of new physics are substantially higher with respect to the “minimal” case, where the Yukawa couplings are assumed to be the only breaking sources of both symmetries [41].

If $\tan \beta$ is large, the inclusion of flavor-diagonal phases has interesting effects also in flavour-changing processes. The main consequences, derived in a model independent manner, can be summarized as follows [39]: (i) extra CPV can only arise from flavor diagonal CPV sources in the UV theory; (ii) the extra CP phases in $B_s - \bar{B}_s$ mixing provide an upper bound on the amount of CPV in $B_d - \bar{B}_d$ mixing; (iii) if operators containing right-handed light quarks are sub-dominant then the extra CPV is equal in the two systems, and is negligible in transitions between the first two generations to the third one. Conversely, these operators can break the correlation between CPV in the $B_s$ and $B_d$ systems, and can induce significant new CPV in $\epsilon_K$. 

V. SUPERSYMMETRY

Supersymmetric models provide, in general, new sources of flavor violation, for both the quark and the lepton sectors. The main new sources are the supersymmetry breaking soft mass terms for squarks and sleptons, and the trilinear couplings of a Higgs field with a squark-antisquark, or slepton-antislepton pairs. Let us focus on the squark sector. The new sources of flavor violation are most commonly analyzed in the basis in which the corresponding (down or up) quark mass matrix and the neutral gaugino vertices are diagonal. In this basis, the squark masses are not necessarily flavor-diagonal, and have the form

$$\tilde{q}_M^* (M^2_{\tilde{q}})_{ij} M_N \tilde{q}_N = (\tilde{q}_{L_i}^* \tilde{q}_{R_k}) \left( \begin{array}{cc} (M^2_{\tilde{d}})_{Lij} & A^d_{ij} v_d \\ A^d_{ij} v_d & (M^2_{\tilde{u}})_{Rkl} \end{array} \right) \left( \begin{array}{c} \tilde{q}_{L_j} \\ \tilde{q}_{R_l} \end{array} \right),$$

where $M, N = L, R$ label chirality, and $i, j, k, l = 1, 2, 3$ are generation indices. $(M^2_{\tilde{q}})_{LL}$ and $(M^2_{\tilde{q}})_{RR}$ are the supersymmetry-breaking squark masses-squared. The $A^q$ parameters enter in the trilinear scalar couplings $A^q_{ij} \phi q \tilde{q}_L \tilde{q}_R$, where $\phi$ ($q = u, d$) is the $q$-type Higgs boson and $v_q = \langle \phi_q \rangle$.

In this basis, flavor violation takes place through one or more squark mass insertion. Each mass insertion brings with it a factor of $\delta q_{ij}^{MN} \equiv \left( \frac{M^2_{\tilde{q}}}{\tilde{m}_q^2} \right)_{ij}^{MN}$, where $\tilde{m}_q^2$ is a representative $q$-squark mass scale. Physical processes therefore constrain

$$[(\delta^q_{ij})_{MN}]_{\text{eff}} \sim \max[(\delta^q_{ij})_{MN}, (\delta^q_{ik})_{MP} (\delta^q_{kj})_{PN}, \ldots, (i \leftrightarrow j)].$$

For example,

$$[(\delta^d_{12})_{LR}]_{\text{eff}} \sim \max[A^d_{12} v_d / \tilde{m}_d^2, (M^2_{\tilde{d}})_{L1k} A^d_{k2} v_d / \tilde{m}_d^4, A^d_{1k} v_d(M^2_{\tilde{d}})_{Rk2} / \tilde{m}_d^4, \ldots, (1 \leftrightarrow 2)].$$

Note that the contributions with two or more insertions may be less suppressed than those with only one.

In terms of mass basis parameters, the $(\delta^q_{ij})_{MM}$’s stand for a combination of mass splittings and mixing angles:

$$\langle \delta^q_{ij} \rangle_{MM} = \frac{1}{\tilde{m}_q^2} \sum_{\alpha} (K^q_M)_{i\alpha} (K^q_M)_{j\alpha} \Delta \tilde{m}_{q_{\alpha}}^2,$$

where $K^q_M$ is the mixing matrix in the coupling of the gluino (and similarly for the bino and neutral wino) to $q_L - \tilde{q}_M; \tilde{m}_q^2 = \frac{1}{3} \sum_{\alpha=1}^{3} \tilde{m}_{q_{\alpha}}$ is the average squark mass-squared, and $\Delta \tilde{m}_q^2 = \tilde{m}_{q_{\alpha}}^2 - \tilde{m}_q^2$.
\[ |K_{ik}K_{jk}^*| \ll |K_{ij}K_{jj}^*|, \quad |K_{ik}K_{jk}^*\Delta \hat m_{q_{ij}}^2| \ll |K_{ij}K_{jj}^*\Delta \hat m_{q_{ij}}^2|, \]  

(5.5)

where there is no summation over \( i, j, k \) and where \( \hat m_{q_{ij}}^2 = \hat m_{q_j}^2 - \hat m_{q_i}^2 \). Then, the contribution of the intermediate \( \hat q_k \) can be neglected and, furthermore, to a good approximation, \( K_{il}K_{lj}^* + K_{ij}K_{jj}^* = 0 \).

For these cases, we obtain

\[ (\delta_{ij})_{MM} = \frac{\Delta \hat m_{q_{ij}}^2}{\hat m_{q_j}^2} (K_{qM}^q)_{ij} (K_{qM}^q)_{jj}^*. \]  

(5.6)

It is further useful to use instead of \( \hat m_q \) the mass scale \( \hat m_{q_{ij}}^q = \frac{1}{2}(\hat m_{q_i} + \hat m_{q_j}) \) \[43\]. We also define

\[ \langle \delta_{ij} \rangle = \sqrt{(\delta_{ij})_{LL}(\delta_{ij})_{RR}}. \]  

(5.7)

The new sources of flavor and CP violation contribute to FCNC processes via loop diagrams involving squarks and gluinos (or electroweak gauginos, or higgsinos). If the scale of the soft supersymmetry breaking is below TeV, and if the new flavor violation is of order one, and/or if the phases are of order one, then these contributions could be orders of magnitude above the experimental bounds. Imposing that the supersymmetric contributions do not exceed the phenomenological constraints leads to constraints of the form \( (\delta_{ij})_{MM} \ll 1 \). Such constraints imply that either quasi-degeneracy \( (\Delta \hat m_{q_{ij}}^2 \ll \hat m_{q_{ij}}^2) \) or alignment \( (|K_{qM}^q| \ll 1) \) or a combination of the two mechanisms is at work.

Table IV presents the constraints obtained in Refs. \[44–47\] as appear in \[42\]. Wherever relevant, a phase suppression of order 0.3 in the mixing amplitude is allowed, namely we quote the stronger between the bounds on \( \Re (\delta_{ij}) \) and \( 3\Im (\delta_{ij}) \). The dependence of these bounds on the average squark mass \( \hat m_{\tilde q} \), the ratio \( x \equiv m_{\tilde g}^2/m_{\tilde q}^2 \) as well as the effect of arbitrary strong CP violating phases can be found in \[42\].

For large \( \tan \beta \), some constraints are modified from those in Table IV. For instance, the effects of neutral Higgs exchange in \( B_s \) and \( B_d \) mixing give, for \( \tan \beta = 30 \) and \( x = 1 \) (see \[42, 48\] and references therein for details):

\[ \langle \delta_{13}^d \rangle < 0.01 \cdot \left( \frac{M_{A^0}}{200 \text{ GeV}} \right), \quad \langle \delta_{23}^d \rangle < 0.04 \cdot \left( \frac{M_{A^0}}{200 \text{ GeV}} \right), \]  

(5.8)

where \( M_{A^0} \) denotes the pseudoscalar Higgs mass, and the above bounds scale roughly as \((30/\tan \beta)^2\).

The experimental constraints on the \( (\delta_{ij})_{LR} \) parameters in the quark-squark sector are presented in Table V. The bounds are the same for \( (\delta_{ij})_{LR} \) and \( (\delta_{ij})_{RL} \), except for \( (\delta_{12}^d)_{MN} \), where the bound
\begin{table}
\centering
\begin{tabular}{|c|cc|}
\hline
$q\ ij$ & $\langle \delta_{ij}^q \rangle_{MM}$ & $\langle \delta_{ij}^q \rangle$ \\
\hline
d 12 & 0.03 & 0.002 \\
d 13 & 0.2 & 0.07 \\
d 23 & 0.6 & 0.2 \\
u 12 & 0.1 & 0.008 \\
\hline
\end{tabular}
\caption{The phenomenological upper bounds on $\langle \delta_{ij}^q \rangle_{MM}$ and on $\langle \delta_{ij}^q \rangle$, where $q = u,d$ and $M = L,R$.}
\end{table}

The constraints are given for $m_\tilde{q} = 1$ TeV and $x \equiv m_\tilde{q}^2/m_\tilde{g}^2 = 1$. We assume that the phases could suppress the imaginary parts by a factor $\sim 0.3$. The bound on $\langle \delta_{23}^d \rangle_{RR}$ is about 3 times weaker than that on $\langle \delta_{23}^d \rangle_{LL}$ (given in table). The constraints on $(\delta_{12,13})_{MM}$, $(\delta_{12}^u)_{MM}$ and $(\delta_{23}^d)_{MM}$ are based on, respectively, Refs. \cite{44,45,46}.

\begin{table}
\centering
\begin{tabular}{|c|cc|}
\hline
$q\ ij$ & $\langle \delta_{ij}^q \rangle_{LR}$ \\
\hline
d 12 & $2 \times 10^{-4}$ \\
d 13 & 0.08 \\
d 23 & 0.01 \\
d 11 & $4.7 \times 10^{-6}$ \\
u 11 & $9.3 \times 10^{-6}$ \\
u 12 & 0.02 \\
\hline
\end{tabular}
\caption{The phenomenological upper bounds on chirality-mixing $(\delta_{ij}^q)_{LR}$, where $q = u,d$. The constraints are given for $m_\tilde{q} = 1$ TeV and $x \equiv m_\tilde{q}^2/m_\tilde{g}^2 = 1$. The constraints on $\delta_{12,13}^d$, $\delta_{12}^u$, $\delta_{23}^d$ and $\delta_{11}^u$ are based on, respectively, Refs. \cite{44,45,46,49} (with the relation between the neutron and quark EDMs as in \cite{50}).}
\end{table}

for $MN = LR$ is 10 times weaker. Very strong constraints apply for the phase of $(\delta_{11}^u)_{LR}$ from EDMs. For $x = 4$ and a phase smaller than 0.1, the EDM constraints on $(\delta_{11}^{u,d,\ell})_{LR}$ are weakened by a factor $\sim 6$.

While, in general, the low energy flavor measurements constrain only the combinations of the suppression factors from degeneracy and from alignment, such as Eq. (5.6), an interesting exception occurs when combining the measurements of $K^0-\bar{K}^0$ and $D^0-\bar{D}^0$ mixing to test the first two generation squark doublets. Here, for masses below the TeV scale, some level of degeneracy is unavoidable \cite{16}:

\begin{equation}
\frac{m_{\tilde{q}_2} - m_{\tilde{q}_1}}{m_{\tilde{q}_2} + m_{\tilde{q}_1}} \leq \begin{cases} 
0.034 & \text{maximal phases} \\
0.27 & \text{vanishing phases} 
\end{cases} \tag{5.9}
\end{equation}
The strong constraints in Tables IV and V can be satisfied if the mediation of supersymmetry breaking to the MSSM is MFV. In particular, if at the scale of mediation, the supersymmetry breaking squark masses are universal, and the A-terms vanish or are proportional to the Yukawa couplings, then the model is phenomenologically safe. Indeed, there are several known mechanisms of mediation that are MFV (see, e.g. [51]). In particular, gauge-mediation [52–55], anomaly-mediation [56, 57], and gaugino-mediation [58] are such mechanisms. (The renormalization group flow in the MSSM with generic MFV soft-breaking terms at some high scale has recently been discussed in Ref. [41, 59].) On the other hand, we do not expect gravity-mediation to be MFV and it could provide subdominant, yet observable flavor and CP violating effects [60].

VI. EXTRA DIMENSIONS

Models of extra dimensions come in a large variety and the corresponding phenomenology, including the implications for flavor physics, changes from one extra dimension framework to another. Yet, as in the supersymmetric case, one can classify the new sources of flavor violation which generically arise:

Bulk masses - If the SM fields propagate in the bulk of the extra dimensions they can have bulk, vector-like, masses. These mass terms are of particular importance to flavor physics since they induce fermion localization which may yield hierarchies in the low energy effective couplings. Furthermore, the bulk masses, which define the extra dimension interaction basis, do not need to commute with the Yukawa matrices, and hence might induce contributions to FCNC processes, similarly to the squark soft masses-squared in supersymmetry.

Cutoff, UV physics - Since, generically, higher dimensional field theories are non-renormalizable, they rely on unspecified microscopic dynamics to provide UV completion of the models. Hence, they can be viewed as effective field theories and the impact of the UV physics is expected to be captured by a set of operators suppressed by the, framework dependent, cutoff scale. Without precise knowledge of the short distance dynamics, the additional operators are expected to carry generic flavor structure and contribute to FCNC processes. This is somewhat similar to “gravity mediated” contributions to supersymmetry-breaking soft terms which are generically expected to have an anarchic flavor structure and are suppressed by the Planck scale.

“Brane” localized terms - The extra dimensions have to be compact and typically contain defects and boundaries of smaller dimensions [in order, for example, to yield a chiral low energy four dimension (4D) theory]. These special points might contain different microscopical degrees
of freedom. Therefore, generically, one expects that a different and independent class of higher dimension operators may be localized to this singular region in the extra dimension manifold. (These are commonly denoted ‘brane terms’ even though, in most cases, they have very little to do with string theory). The brane-localized terms can, in principle, be of anarchic flavor structure and provide new flavor and CP violating sources. One important class of such operators are brane kinetic terms: their impact is somewhat similar to that of non-canonical kinetic terms which generically arise in supersymmetric flavor models.

We focus on flavor physics of five dimension (5D) models, with bulk SM fields, since most of the literature focuses on this class. Furthermore, the new flavor structure that arises in 5D models captures most of the known effects of extra-dimension flavor models. Assuming a flat extra dimension, the energy range, $\Lambda_{5D} R$ (where $\Lambda_{5D}$ is the 5D effective cutoff scale and $R$ is the extra dimension radius with the extra dimension coordinate $y \in (0, \pi R)$), for which the 5D effective field theory holds, can be estimated as follows. Since gauge couplings in extra dimensional theories are dimensional, i.e. $\alpha_{5D}$ has mass dimension $-1$, a rough guess (which is confirmed, up to order one corrections, by various NDA methods) is $\Lambda_{5D} \sim 4\pi/\alpha_{5D}$. Matching this 5D gauge coupling to a 4D coupling of the SM, at leading order, $1/g^2 = \pi R/g_{5D}^2$, we obtain

$$\Lambda_{5D} R \sim \frac{4}{\alpha} \sim 30.$$  \hspace{1cm} (6.1)

Generically, the mass of the lightest Kaluza-Klein (KK) states, $M_{KK}$, is of $\mathcal{O}(R^{-1})$. If the extra dimension theory is linked to the solution of the hierarchy problem and/or directly accessible to near future experiments, then $R^{-1} = \mathcal{O}(\text{TeV})$. This implies an upper bound on the 5D cutoff:

$$\Lambda_{5D} \lesssim 10^2 \text{TeV} \ll \Lambda_K \sim 2 \times 10^5 \text{TeV},$$  \hspace{1cm} (6.2)

where $\Lambda_K$ is the scale required to suppress the generic contributions to $\epsilon_K$, discussed above (see Table 1).

The above discussion ignores the possibility of splitting the fermions in the extra dimension. In split fermion models different bulk masses are assigned to different generations. Consequently fermions are localized and separated in the bulk of the extra dimension in a manner which may successfully address the SM flavor puzzle [62]. Separation in the extra dimension may suppress the contributions to $\epsilon_K$ from the higher-dimension cut-off induced operators. As shown in Table 1, the most dangerous operator is

$$O^4_K = \frac{1}{\Lambda_{5D}^2} (\bar{s}_L d_R) (\bar{s}_R d_L).$$  \hspace{1cm} (6.3)
This operator contains $s$ and $d$ fields of both chiralities. As a result, in a large class of split fermion models, the overlap suppression would be similar to that accounting for the smallness of the down and strange 4D Yukawa couplings. Integrating over the 5D profiles of the four quarks, this may yield a suppression factor of $O(m_d m_s/v^2) \sim 10^{-9}$. Together with the naive scale suppression, $1/\Lambda_{5D}^2$, the coefficient of $O_K^4$ can be sufficiently suppressed to be consistent with the experimental bound.

In the absence of large brane kinetic terms (BKTs), however, fermion localization generates order one non-universal couplings to the gauge KK fields [63]. (See e.g. [64] and references therein. The case with large BKTs is similar to the warped case discussed below.) The fact that the bulk masses are, generically, not aligned with the 5D Yukawa couplings implies that KK gluon exchange processes induce, among others, the following operator in the low energy theory:

$$\left[(D_L)_{12}/(6M_{KK}^2)\right] (\bar{s}_L d_L)^2,$$

where $(D_L)_{12} \sim \lambda$ is the left-handed down-quark rotation matrix from the 5D interaction basis to the mass basis. This structure provides only a mild suppression to the resulting operator. It implies that to satisfy the $\epsilon_K$ constraint the KK and the inverse compactification scales have to be above $10^3$ TeV, beyond the direct reach of near future experiments, and too high to be linked to a solution of the hierarchy problem. This problem can be solved by tuning the 5D flavor parameters and imposing appropriate 5D flavor symmetries to make the tuning stable. Once the 5D bulk masses are aligned with the 5D Yukawa matrices the KK gauge contributions would vanish and such a configuration is radiatively stable.

The warped extra-dimension [Randall-Sundrum (RS)] framework [65] provides a solution to the hierarchy problem. Moreover, with SM fermions propagating in the bulk, both the SM and the NP flavor puzzles can be addressed. The light fermions could be localized away from the TeV brane [66], where the Higgs is localized. Such a configuration can generate the observed Yukawa hierarchy and, at the same time, ensure that higher-dimensional operators are suppressed by a high cutoff scale associated with the location of the light fermions in the extra dimension [67]. Furthermore, since the KK states are localized near the TeV brane, the couplings between the SM quarks and the gauge KK fields exhibit the hierarchical structure associated with SM masses and CKM mixings. This hierarchy in the couplings provides an extra protection against non-standard flavor-violating effects [68] denoted as RS-GIM mechanism [69] (see also [70]). It is interesting to note that an analogous mechanism is at work in models with strong dynamics at the TeV scale, with large anomalous dimension and partial compositeness [71]. The link with strongly-interacting models in indeed motivated by the AdS/CFT correspondence [72], which implies that the above 5D framework is the dual description of 4D composite Higgs models [73].
Concerning the quark zero modes, the flavor structure of the above models as well as the phenomenology can be captured by using the following simple rules [69, 74, 75]. In the 5D interaction basis, where the bulk masses \( kC_{ij} \) are diagonal (\( x = Q, U, D; i, j = 1, 2, 3; k \) is the AdS curvature), the value \( f_{x_i} \) of the profile of the quark zero modes is given by

\[
f_{x_i}^2 = (1 - 2c_{x_i})/(1 - \epsilon^{1 - 2c_{x_i}}) .
\]  
(6.4)

Here \( c_{x_i} \) are the eigenvalues of the \( C_x \) matrices, \( \epsilon = \exp[-\xi], \xi = \log[M_{Pl}/\text{TeV}] \), and \( M_{Pl} \) is the reduced Planck mass. If \( c_{x_i} < 1/2 \), then \( f_{x_i} \) is exponentially suppressed. Hence, order one variations in the 5D masses yield large hierarchies in the 4D flavor parameters. We consider the cases where the Higgs VEV either propagates in the bulk [76] or is localized on the IR brane. For a bulk Higgs case, the profile is given by

\[
\tilde{v}(\beta, z) \approx v\sqrt{k(1 + \beta)}\bar{z}^{2 + \beta}/\epsilon, \quad \bar{z} \in (\epsilon, 1) \quad (\bar{z} = 1 \text{ on the IR brane}), \quad \beta \geq 0 .
\]

The \( \beta = 0 \) case describes a Higgs maximally-spread into the bulk (saturating the AdS stability bound [77]). The relevant part of the effective 4D Lagrangian, which involves the zero modes and the first KK gauge states can be approximated by [69, 74]

\[
\mathcal{L}^{4D} \supset (Y_{5D})_{ij}^u \phi^u_0 \bar{Q}_i f_{U,D} \tilde{Q}_i, f_{U,D} r_{00}^\phi(\beta, c_{Q_1}, c_{U,D}j) + g_s G^1 x_i^x_i [f_{x_i}^2 + g^0_0(c_{x_i}) - 1/\xi] ,
\]  
(6.5)

where \( \phi^u_0 = \tilde{\phi}, \phi \), \( g_s \) stands for a generic effective gauge coupling and summation over \( i,j \) is implied. The correction for the couplings from the case of fully IR-localized KK and Higgs states is given by the functions \( r_{00}^\phi [74] \) and \( r_{00}^g [78, 79] \):

\[
r_{00}^\phi(\beta, c_L, c_R) \approx \frac{\sqrt{2(1 + \beta)}}{2 + \beta - c_L - c_R}, \quad r_{00}^g(c) \approx \frac{\sqrt{2}}{J_1(x_1)} \frac{0.7}{6 - 4c} \left(1 + e^{c/2}\right) ,
\]  
(6.6)

where \( r_{00}^\phi(\beta, c_L, c_R) = 1 \) for brane-localized Higgs and \( x_1 \approx 2.4 \) is the first root of the Bessel function, \( J_0(x_1) = 0 \).

In Table VI we present an example of a set of \( f_{x_i} \)-values that, starting from anarchical 5D Yukawa couplings, reproduce the correct hierarchy of the flavor parameters. We assume, for simplicity, an IR localized Higgs. The values depend on two input parameters: \( f_{U,3} \), which has been determined assuming a maximally localized \( t_R \) (\( c_{u_3} = -0.5 \)), and \( y_{5D} \), the overall scale of the 5D Yukawa couplings in units of \( k \), which has been fixed to its maximal value assuming three KK states. On general grounds, the value of \( y_{5D} \) is bounded from above, as a function of the number of KK levels, by the requirement that Yukawa interactions are perturbative below the cutoff of the theory, \( \Lambda_{5D} \sim N_{KK} k \), and it is bounded from below in order to account for the large top mass. Hence the following range for \( y_{5D} \) is obtained (see e.g. [76, 81]):

\[
\frac{1}{2} \lesssim y_{5D} \lesssim \frac{2\pi}{N_{KK}} \quad \text{for brane Higgs}; \quad \frac{1}{2} \lesssim y_{5D} \lesssim \frac{4\pi}{\sqrt{N_{KK}}} \quad \text{for bulk Higgs},
\]  
(6.7)
where we use the rescaling $y_{5D} \rightarrow y_{5D} \sqrt{1 + \beta}$, which produces the correct $\beta \rightarrow \infty$ limit \cite{ref6} and avoids subtleties in the $\beta = 0$ case.

With anarchical 5D Yukawa matrices, an RS residual little CP problem remains \cite{ref80}: Too large contributions to the neutron electric dipole moment (EDM) \cite{ref39}, and sizable chirally enhanced contributions to $\epsilon_K$ \cite{ref2, ref78, ref24, ref25} are predicted. The RS leading contribution to $\epsilon_K$ is generated by a tree-level KK-gluon exchange which leads to an effective coupling for the chirality-flipping operator in (6.3) of the type \cite{ref78, ref24, ref25}

$$C_4^K \simeq \frac{g_{sss}^2}{M_{KK}^2} f_{Q_2} f_{Q_1} f_{d_2} f_{d_1} r_{100}^g(c_{Q_2}) r_{00}^g(c_{d_2}) \sim \frac{g_{sss}^2}{M_{KK}^2 (vy_{5D})^2} \frac{m_d m_s}{M_{KK}^2} r_{100}^H(c_{Q_2}) r_{00}^H(c_{d_2}).$$

(6.8)

The final expression is independent of the $f_{x_i}$, so the bound in Table I can be translated into constraints in the $y_{5D} - M_{KK}$ plane. The analogous effects in the $D$ and $B$ systems yield numerically weaker bounds. Another class of contributions, which involves only left-handed quarks, is also important to constrain the $f_Q - M_{KK}$ parameter space.
In Table VII we summarize the resulting constraints. For the purpose of a quantitative analysis we set \( g_{s} = 3 \), as obtained by matching to the 4D coupling at one-loop [76] (for the impact of a smaller RS volume see [84]). The constraints related to CPV correspond to maximal phases, and are subject to the requirement that the RS contributions are smaller than 30% (60%) of the SM contributions \( \epsilon \) in the \( B_d \) \((K)\) system. The analytical expressions in the table have roughly a 10% accuracy over the relevant range of parameters. Contributions from scalar exchange, either Higgs [81] or radion [86], are not included since these are more model dependent and known to be weaker [87] in the brane localized Higgs case.

Constraints from \( \epsilon' / \epsilon_K \) have a different parameter dependence than the \( \epsilon_K \) constraints. Explicitly, for \( \beta = 0 \), the \( \epsilon' / \epsilon_K \) constraint reads \( M_{G}^{\text{min}} = 1.2 y_{5D} \) TeV. When combined with the \( \epsilon \) constraint, we find \( M_{G}^{\text{min}} = 5.5 \) TeV with a corresponding \( y_{5D}^{\text{min}} = 4.5 \) [74].

The constraints summarized in Table VII and the contributions to the neutron EDM which generically require \( M_{KK} > \mathcal{O}(20 \text{ TeV}) \) [83] are a clear manifestation of the RS little CP problem. The problem can be amended by various alignment mechanisms [79, 80, 88]. In this case the bounds from the up sector, especially from CPV in the \( D \) system [16, 47], become important. Constraint from \( \Delta F = 1 \) processes (in either the down sector [39, 84] or \( t \to cZ \) [89]) are not included here, since they are weaker and, furthermore, these contributions can be suppressed (see [84]) due to incorporation of a custodial symmetry [90].

VII. FUTURE PROSPECTS

The new physics flavor puzzle is the question of why, and in what way, the flavor structure of TeV-scale new physics is non-generic. Indeed, the flavor predictions of most new physics models are not a consequence of their generic features but rather of the special structures that are imposed specifically to satisfy the existing severe flavor bounds. Therefore, flavor physics is a powerful indirect probe of new physics. We hope that new physics not far above the weak scale will be discovered at the LHC. A major issue will then be to understand its flavor structure. While it is not easy to directly probe this flavor structure at high energy, a lot can be learned from low energy flavor physics.

The precision with which we can probe high scale physics in flavor physics experiments is often limited by theoretical uncertainties. Moreover, in case of theoretically clean observables, the sensitivity to the new-physics scale increases slowly with the statistics of the experiment. Thus, the important questions in view of future experiments are the following:
1. What are the expected deviations from the SM predictions induced by new physics at the TeV scale?

2. Which observables are not limited by theoretical uncertainties?

3. In which case we can expect a substantial improvement on the experimental side?

4. What will the measurements teach us if deviations from the SM are [not] seen?

These questions have been analysed in a series of recent works (see e.g. [46, 49, 91–93]) and the main conclusions can be summarized as follows:

1. The expected deviations from the SM predictions induced by new physics at the TeV scale with generic flavor structure are already ruled out by many orders of magnitudes. On general grounds, we can expect any size of deviation below the current bounds. In the most pessimistic frameworks, such as MFV, the typical size of the deviations is at the few % level in FCNC amplitudes.

2. The theoretical limitations are highly process dependent. In most multi-hadron final states the calculation of decay amplitudes are already limited by theoretical uncertainties. However, several channels involving leptons in the final state, and selected time-dependent asymmetries, have a theoretical errors well below the current experimental sensitivity.

3. On the experimental side there are good prospect of improvements. As summarized in Table VIII, one order of magnitude improvements in several clean $B_{s,d}$, $D$, and $K$ observables are possible within a few years. Moreover, improvements of several orders of magnitudes are expected in top decays, which will be explored for the first time in great detail at the LHC.

4. There is no doubt that new low-energy flavor data will be complementary with the high-$p_T$ part of the LHC program. As illustrated in the previous sections, the synergy of both data sets can teach us a lot about the new physics at the TeV scale.

While some improvements can still be expected from running experiments, in particular from CDF and D0 at Fermilab, a substantial step forward can be achieved only with the new dedicated facilities. At the LHC, which has just started to operate, the LHCb experiment is expected to collect about 10 fb$^{-1}$ of data by 2015. Beyond that, an LHCb upgrade is planned with 10 times larger luminosity. At CERN the NA62 experiment is expected to start in 2012, with the main goal of collecting $\mathcal{O}(100)$ events of the rare decay $K^+ \rightarrow \pi^+\nu\bar{\nu}$ by 2014. Dedicated rare $K$ experiments
| Observable | SM prediction | Theory error | Present result | Future error | Future Facility |
|------------|---------------|--------------|----------------|--------------|----------------|
| $|V_{us}|$   | $[K \to \pi \ell \nu]$ | input         | $0.5\% \to 0.1\%_{\text{Latt}}$ | $0.2246 \pm 0.0012$ | $0.1\%$ | $K$ factory |
| $|V_{cb}|$   | $[B \to X_s \ell \nu]$ | input         | $1\%$          | $(41.54 \pm 0.73) \times 10^{-3}$ | $1\%$ | Super-B     |
| $|V_{ub}|$   | $[B \to \pi \ell \nu]$ | input         | $10\% \to 5\%_{\text{Latt}}$ | $(3.38 \pm 0.36) \times 10^{-3}$ | $4\%$ | Super-B     |
| $\gamma$   | $[B \to DK]$ | input         | $< 1^\circ$     | $(70^{+27}_{-30})^\circ$ | $3^\circ$ | LHCb        |
| $S_{B_d \to \psi K}$ | $\sin(2\beta)$ | $\lesssim 0.01$ | $0.671 \pm 0.023$ | $0.01$ | LHCb |
| $S_{B_s \to \psi}$ | $0.036$ | $\lesssim 0.01$ | $0.81^{+0.12}_{-0.32}$ | $0.01$ | LHCb |
| $S_{B_d \to \phi K}$ | $\sin(2\beta)$ | $\lesssim 0.05$ | $0.44 \pm 0.18$ | $0.1$ | LHCb |
| $S_{B_s \to \phi \phi}$ | $0.036$ | $\lesssim 0.05$ | $-0.016 \pm 0.22$ | $0.05$ | LHCb |
| $S_{B_d \to K^* \phi \gamma}$ | few $\times 0.01$ | $0.01$ | $-0.16 \pm 0.22$ | $0.05$ | LHCb |
| $A_{S_{B_d \to \phi \gamma}}$ | few $\times 0.01$ | $0.01$ | $-0.16 \pm 0.22$ | $0.05$ | LHCb |
| $A_{B_d}^d$ | $-5 \times 10^{-4}$ | $10^{-4}$ | $(5.8 \pm 3.4) \times 10^{-3}$ | $10^{-3}$ | LHCb |
| $A_{S_{B_d}}^d$ | $2 \times 10^{-5}$ | $< 10^{-5}$ | $(1.6 \pm 8.5) \times 10^{-3}$ | $10^{-3}$ | LHCb |
| $AC_{\gamma}(b \to s \gamma)$ | $< 0.01$ | $< 0.01$ | $-0.012 \pm 0.028$ | $0.005$ | Super-B |
| $B(B \to \tau \nu)$ | $1 \times 10^{-4}$ | $20\% \to 5\%_{\text{Latt}}$ | $(1.73 \pm 0.35) \times 10^{-4}$ | $5\%$ | Super-B |
| $B(B \to \mu \nu)$ | $4 \times 10^{-7}$ | $20\% \to 5\%_{\text{Latt}}$ | $< 1.3 \times 10^{-6}$ | $6\%$ | Super-B |
| $B(B_s \to \mu^+ \mu^-)$ | $3 \times 10^{-9}$ | $20\% \to 5\%_{\text{Latt}}$ | $< 5 \times 10^{-8}$ | $10\%$ | LHCb |
| $B(B_d \to \mu^+ \mu^-)$ | $1 \times 10^{-10}$ | $20\% \to 5\%_{\text{Latt}}$ | $< 1.5 \times 10^{-8}$ | $[?]$ | LHCb |
| $A_{FB}(B \to K^{\ast} \mu^+ \mu^-)_{q\bar{q}}$ | $0$ | $0.05$ | $(0.2 \pm 0.2)$ | $0.05$ | LHCb |
| $B \to K \nu \bar{\nu}$ | $4 \times 10^{-6}$ | $20\% \to 10\%_{\text{Latt}}$ | $< 1.4 \times 10^{-5}$ | $20\%$ | Super-B |
| $|q/p|_{D\text{-mixing}}$ | $1$ | $< 10^{-3}$ | $(0.86^{+0.18}_{-0.15})$ | $0.03$ | Super-B |
| $\phi_D$ | $0$ | $< 10^{-3}$ | $(9.6^{+8.4}_{-0.5})^\circ$ | $2^\circ$ | Super-B |
| $B(K^+ \to \pi^+ \nu \bar{\nu})$ | $8.5 \times 10^{-11}$ | $8\%$ | $(1.73^{+1.5}_{-0.75}) \times 10^{-10}$ | $10\%$ | $K$ factory |
| $B(K_L \to \rho^0 \nu \bar{\nu})$ | $2.6 \times 10^{-11}$ | $10\%$ | $< 2.6 \times 10^{-8}$ | $[?]$ | $K$ factory |
| $R(c/\mu)(K \to \pi \ell \nu)$ | $2.477 \times 10^{-5}$ | $0.04\%$ | $(2.498 \pm 0.014) \times 10^{-5}$ | $0.1\%$ | $K$ factory |
| $B(t \to c Z, \gamma)$ | $\mathcal{O}(10^{-13})$ | $\mathcal{O}(10^{-13})$ | $< 0.6 \times 10^{-2}$ | $\mathcal{O}(10^{-5})$ | LHC (100 fb$^{-1}$) |

TABLE VIII: Status and prospects of selected $B_{s,d}$, $D$, $K$ and $t$ observables (based on information from Ref. [17, 31, 22]). In the third column “Latt” refer to improvements in Lattice QCD expected in the next 5 years. In the fourth column the bounds are 90% CL. The errors in the fifth column refer to 10 fb$^{-1}$ at LHCb, 50 ab$^{-1}$ at Super-B, and two years at NA62 (“K factory”). In the third and fifth column the errors followed by “%” are relative errors, while the others are absolute errors. For entries marked “[?]” we have not found a reliable estimate of the future experimental prospects.
are planned also at JParc and at Fermilab. The project to upgrade KEK-B into a Super-$B$ factory, with a luminosity exceeding $5 \times 10^{35}/\text{cm}^2\text{s}^{-1}$ has just started, and an even more ambitious super-$B$ project is currently under study in Italy.

In Table VIII, based on data from Ref. [17, 18, 46, 91, 92], we show the possible improvements for a series of particularly significant $B_{s,d}$, $D$, $K$ and $t$ observables, most of which have already been discussed in the previous sections. The future improvements refer to $10 \text{fb}^{-1}$ at LHCb, $50 \text{ab}^{-1}$ at Super-$B$, and two years of nominal data taking at NA62. The table is not comprehensive and the entries in the last two columns necessarily have significant uncertainties. However, it illustrates two important points: i) there are still several clean observables where we can expect significant improvements in the near future; ii) progress in this field requires the combined efforts of different experimental facilities.

To conclude, we stress that the future of flavor physics is promising:

- The technology to collect much more flavor data exists.
- There are still several measurements which are not theory limited.
- Most well-motivated TeV-scale extensions of the Standard Model predict deviations within future experimental sensitivities and above the SM theoretical uncertainties.

The combination of direct discoveries at the LHC and a pattern of deviations in flavor machines is likely to solve the new physics flavor puzzle, to indirectly probe physics at a scale much higher than the TeV scale, and, perhaps to make progress in understanding the standard model flavor puzzle.

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