SUSY beyond minimal flavour violation

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Abstract. We review aspects of the phenomenology of the MSSM with non-minimal flavour violation, including a discussion of important constraints and the sensitivity to fundamental scales.

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1 Introduction

The Standard Model (SM) of particle physics provides an economical description of thousands of observables in particle physics (and presumably most everyday phenomena) based on the principle of spontaneously broken gauge invariance. Nevertheless, there are phenomena which do not find a straightforward explanation within the SM, notably dark matter, neutrino masses, and the gravitational force. The latter two, together with the approximate unification of the three running couplings in the ultraviolet, provide three independent hints at new dynamics at scales $M_X \sim \mathcal{O}(10^{15} - 10^{19})$ GeV. Explaining and stabilizing the ensuing hierarchy $M_W \ll M_X$ provideds some of the principal motivation for supersymmetry, the eponym of this conference. SUSY necessarily entails an enlargement of the particle spectrum (including species suitable to make up the observed dark matter), naturally improves the unification of couplings, and makes the hierarchy stable against quantum corrections.

The LHC is designed to directly probe the TeV scale and clarify the mechanism of electroweak symmetry breaking. If supersymmetry is involved, ATLAS and CMS will likely detect at least part of the particle spectrum directly. On the other hand, the examples given above demonstrate that low-energy observables exist that can probe fundamental scales, including those beyond the "energy frontier". This happens, for instance, if the fundamental physics violates accidental symmetries of the low-energy theory, i.e. lepton flavour in the case of neutrino oscillations in the context of the seesaw mechanism. SM contributions to hadronic flavour transitions are constrained by the chiral structure of weak interactions and suppressed by the weak scale and the Cabibbo-Kobayashi-Maskawa (CKM) hierarchy, and, in the case of flavour-changing neutral current (FCNC) processes, also by loop factors. Hence these processes, too, provide very sensitive indirect probes.

The remainder of this document is organized as follows. We first review the Lagrangian of the Minimal Supersymmetric Standard Model (MSSM) and the anatomy of the effective vertices. Subsequently we consider some prominent B-physics observables where sparticles may be involved, as well as attempts to combine various measurements to get detailed information on the Lagrangian. We close with patterns predicted in specific scenarios of grand unification. Throughout, the presentation aims at putting self-containedness over exhaustiveness.

2 SUSY flavour violation

2.1 Lagrangian

The minimal supersymmetric standard model (MSSM) [4,5,6] pairs each SM fermion multiplet with scalar partners into a chiral superfield and the gauge bosons with gauginos into vector supermultiplets. The SM Higgs doublet is replaced by two chiral multiplets H_u, H_d (as also required by supersymmetry), each containing a scalar higgs doublet. Imposing R-parity to suppress baryon and lepton number violation, the most general renormalizable superpotential is then fixed by the SM Yukawa couplings, up to one higgs mass parameter (the μ-parameter). In particular, the quartic Higgs couplings are fixed in terms of the gauge couplings by the supersymmetry, and both supersymmetry and the electroweak symmetry are unbroken and guaranteed to remain so at all orders of perturbation theory [1,2]. The phenomenologically required supersymmetry breaking can be introduced explicitly as long as it is soft, i.e. does not introduce quadratic divergences. (In particular, the soft-breaking terms themselves are only logarithmically sensitive to heavy scales

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such as the seesaw scale.) For suitable choices of parameters, the electroweak symmetry is also broken while the hierarchy $M_W \sim M_{\rm SUSY} \ll M_X$ is stabilized, where M_X denotes one of the large scales $M_{\rm seesaw}$, $M_{\rm GUT}$, $M_{\rm Pl}$. More fundamentally, the relevant extra renormalizable terms $\mathcal{L}_{\rm soft}$ in the scalar potential may be due to spontaneous SUSY breaking, which, if dynamical, can naturally generate the large hierarchy $M_{\rm SUSY} \ll M_X$ by dimensional transmutation [3].

Potentially, $\mathcal{L}_{\rm soft}$ contains a large number of flavour- and CP-violating parameters. A general parameterization of the soft-breaking terms is given by a set of trilinear scalar couplings T_U, T_E, T_D (counterparts of the trilinear (Yukawa) couplings in the superpotential), by explicit (generation-dependent) scalar mass terms for the 5 types of chiral multiplets, explicit gaugino masses and three parameters in the Higgs sector. With the exception of the latter two sets, the physical parameters in $\mathcal{L}_{\rm soft}$ are fully determined by the 6×6 hermitian sfermion mass matrices in the so-called super-CKM basis [7,8] (of superfields),

$$M^2_{(\tilde u,\tilde d,\tilde e)} = \begin{pmatrix} M^2_{(\tilde u,\tilde d,\tilde e)LL} & M^2_{(\tilde u,\tilde d,\tilde e)LR} \\ \left(M^2_{(\tilde u,\tilde d,\tilde e)LR}\right)^\dagger & M^2_{(\tilde u,\tilde d,\tilde e)RR} \end{pmatrix}, \qquad (1)$$

together with the 3×3 sneutrino mass matrix $M^2_{\tilde\nu}$. Here,

$$M^2_{\tilde d\,LL} = \hat m^2_Q + m^2_d + D_{d\,LL} \qquad (2)$$
$$M^2_{\tilde u\,LL} = V_{\rm CKM} \hat m^2_Q V^\dagger_{\rm CKM} + m^2_u + D_{u\,LL} \qquad (3)$$
$$M^2_{\tilde d\,RR} = \hat m^2_{\tilde d} + m^2_d + D_{d\,RR} \qquad (4)$$
$$M^2_{\tilde u\,RR} = \hat m^2_{\tilde u} + m^2_u + D_{u\,RR} \qquad (5)$$
$$M^2_{\tilde d\,LR} = v_1 \hat T_D - \mu^* m_d \tan\beta \qquad (6)$$
$$M^2_{\tilde u\,LR} = v_2 \hat T_U - \mu^* m_u \cot\beta \qquad (7)$$
$$D_{f\,LL,RR} = \cos 2\beta\, M_Z^2 (T^3_f - Q_f \sin^2\theta_W)\mathbf{1}_{3\times 3}, \qquad (8)$$

with the charged-slepton mass matrices following via the substitutions $Q \to L$, $d \to e$. Note that the LR masses are proportional to the electroweak scale, hence are suppressed by $v/M_{\rm SUSY}$ in the limit of a large SUSY-breaking scale. The super-CKM basis is defined by requiring diagonal Yukawa couplings and the CKM matrix being in its four-parameter standard form. The neutral fermion-sfermion-gaugino and fermion-sfermion-higgsino couplings are then flavour-diagonal while the charged couplings are governed by the CKM matrix. However, the matrices entering (2)–(7) are in general nondiagonal, subject only to certain hermiticity conditions. The resulting flavour violation is conveniently parameterized in terms of parameters

$$\left(\delta^{u,d,e,\nu}_{ij}\right)_{AB} \equiv \frac{\left[\left(\mathcal{M}_{\tilde u,\tilde d,\tilde e,\tilde\nu}\right)_{AB}\right]_{ij}}{M^2} \qquad (9)$$

where M is a mass scale of order the sfermion mass eigenvalues. (A popular flavour-dependent choice is to use $M^2 = \sqrt{[(\mathcal{M}^2_f)_{XX}]_{ii}[(\mathcal{M}^2)_{YY}]_{jj}}$ in $(\delta^{\tilde f}_{ij})_{AB}$.) When the δ parameters are small, it is possible to treat them as perturbations [8], however caution must be taken to expand to sufficiently high orders to account for all leading effects (see also below).

From the large number of flavour-violating parameters, it is evident that the MSSM generally entails deviations from the SM predictions for flavour-violating processes. This opens the possibility to either observe supersymmetry indirectly or to constrain its parameters (flavour-violating as well as flavour-conserving ones). One virtue of flavour-violating processes is the large number of observables and the availability of theoretical tools for rather precise predictions for many of them. On the other hand, for generic $\delta \sim 1$ and TeV-scale sparticle masses, experimental bounds on FCNC processed such as $B \to X_s\gamma$ are violated (the "SUSY flavour problem"). However, the structure of $\mathcal{L}_{\rm soft}$ is intimately tied to the unknown mechanism of SUSY breaking. For instance, gauge-mediation models [9,10, 11] have little trouble in satisfying the low-energy constraints from flavour physics because the SUSY breaking is tranferred by flavour-blind gauge interactions at relatively low scales.

This spectrum of possibilities is exciting: From a phenomenological point of view, flavour violation provides, as does the sparticle spectrum, nontrivial constraints on dynamics at fundamental scales.

2.2 Effective vertices

In view of the small masses of B, D, and light mesons compared to the weak and SUSY scales, the appropriate tool to separate the heavy scales from the low-energy QCD effects and curb large logarithms is the effective weak Hamiltonian (see [12] for a review of the formalism)

$$\mathcal{H}_{\rm eff} = \sum C_i(\mu) Q_i(\mu), \qquad (10)$$

where Q_i are local operators constructed from the SM fields, C_i are Wilson coefficients encapsulating the effects of the sparticles, and μ is a renormalization (factorization) scale. Integrating out the superpartners at the one-loop level, the operators up to dimension six are determined by penguin and box graphs.

2.2.1 $\Delta F = 2$ (mixing)

The effective $\Delta F = 2$ hamiltonian relevant to meson-antimeson oscillations is solely due to box diagrams. A complete operator basis is given by [13]

$$Q_1^{\rm VLL} = (\bar s^a \gamma_\mu P_L b^a)(\bar s^b \gamma^\mu P_L b^b), \qquad (11)$$
$$Q_1^{\rm LR} = (\bar s^a \gamma_\mu P_L b^a)(\bar s^b \gamma^\mu P_R b^b), \qquad (12)$$
$$Q_2^{\rm LR} = (\bar s^a P_L b^a)(\bar s^b P_R b^b), \qquad (13)$$
$$Q_1^{\rm SLL} = (\bar s^a P_L b^a)(\bar s^b P_L b^b), \qquad (14)$$
$$Q_2^{\rm SLL} = -(\bar s^a \sigma_{\mu\nu} P_L b^a)(\bar s^b \sigma^{\mu\nu} P_L b^b), \qquad (15)$$

(a, b colour indices), plus operators arising from replacing $P_L \leftrightarrow P_R$ in $Q_1^{\rm VLL}$ and $Q_{1,2}^{\rm SLL}$. Only $Q_1^{\rm VLL}$ is generated in the SM (to good approximation). Supersymmetric contributions have been computed in [14,15,16,17,18,19]. For illustration, the SUSY gluino-squark contributions in the mass insertion approximation (expansion in δ's to lowest order) read [15]

$$
\begin{aligned}
C_1^{\rm VLL} &= \epsilon[24 x f_6(x) + 66 \tilde f_6(x)](\delta^d_{sb})^2_{LL}, && (16)\\
C_1^{\rm LR} &= \epsilon[(-12 x f_6(x) - 60 \tilde f_6(x))(\delta^d_{sb})_{LL}(\delta^d_{sb})_{RR} \\
&\quad + 90 \tilde f_6(x)(\delta^d_{sb})_{LR}(\delta^d_{sb})_{RL}], && (17)\\
C_2^{\rm LR} &= \epsilon[(504 x f_6(x) - 72 \tilde f_6(x))(\delta^d_{sb})_{LL}(\delta^d_{sb})_{RR} \\
&\quad - 132 \tilde f_6(x)(\delta^d_{sb})_{LR}(\delta^d_{sb})_{RL}], && (18)\\
C_1^{\rm SLL} &= \epsilon\, 222 x f_6(x)(\delta^d_{sb})^2_{RL}, && (19)\\
C_2^{\rm SLL} &= \epsilon(-9/2) x f_6(x)(\delta^d_{sb})^2_{RL}, && (20)
\end{aligned}
$$

where $\epsilon = -\alpha_s^2/(216\, m_{\tilde q}^2)$ and three more coefficients $C_1^{\rm VRR}$, $C_1^{\rm SRR}$, $C_1^{\rm SRR}$ obtained from $C_1^{\rm VLL}$, $C_1^{\rm SLL}$, $C_1^{\rm SLL}$ via $L \leftrightarrow R$. Here $x = m_{\tilde g}^2/m_{\tilde q}^2$ and $f_6(x)$, $\tilde f_6(x)$ are dimensionless loop functions.

2.2.2 $\Delta F = 1$ (decays)

QCD-penguin graphs contribute only to $\Delta F = 1$ transitions via the operators

$$Q_{3,4} = \sum_q (\bar s_L b_L)\,(\bar q_L q_L), \qquad (21)$$

$$Q_{5,6} = \sum_q (\bar s_L b_L)\,(\bar q_R q_R), \qquad (22)$$

$$Q_{8g} = \bar s_L b_R\, G, \qquad (23)$$

together with operators Q_i' which arise from the Q_i by changing the chiralities of all quarks. Here we have taken the case of $b \to s$ transitions as an example, with obvious replacements for $b \to d$ and $d \to s$ transitions. We have also suppressed colour and part of the Dirac structure. (G is the gluon field strength.) Note that the QCD penguins $Q^{(\prime)}_{3\ldots6}$ involve b and s quarks of like chiralities, while the chromomagnetic penguins $Q^{(\prime)}_{8g}$ involve a chirality flip. This engenders specific patterns of sensitivity of their coefficients to the parameters δ, e.g. $(\delta^d_{sb})_{LR}$ in the case of $C_{7\gamma}$ and C_{8g}.

Photon and Z penguins contribute to operators

$$Q_{9,10} = \sum_q \frac{3 e_q}{2} (\bar s_L b_L)(\bar q_L q_L), \qquad (24)$$

$$Q_{7,8} = \sum_q \frac{3 e_q}{2} (\bar s_L b_L)(\bar q_R q_R), \qquad (25)$$

$$Q_{7\gamma} = \bar s_L\, b_R F, \qquad (26)$$

$$Q_{9V}\,(Q_{10A}) = (\bar s_L b_L)(\bar l_L \bar l_L \pm \bar l_R l_R) \qquad (27)$$

(and their mirror images Q_i'; F is the electromagnetic field strength), which also have analogs contributing to $l_i \to l_j$ transitions. Of the penguin operators, $Q_{3\ldots10}$, $Q_{7\gamma}$, Q_{8g}, Q_{9V}, and Q_{10A} receive significant SM contributions, which are unimportant for the rest. Identifying contributions from any of the primed operators would constitute a clear signal of new physics.

All penguin operators except the magnetic ones also receive contributions from box diagrams; moreover, boxes can contribute to generic four-fermion operators

$$Q = (\bar s_{L,R} b_{L,R})(\bar q_{L,R} q_{L,R}).$$

Note that, in general, these are not suppressed with respect to penguin diagrams [20,15], as all SUSY penguins as well as boxes decouple as M_W^2/M^2 for M large. (This is true even for the Z-penguin, where the bsZ vertex is dimension-four but carries a hidden v^2/M^2 suppression in its coefficient [68].)

Finally, Higgs penguins generating operators such as $\sum_q y_q (\bar s_L b_R)(\bar q_L q_R)$ are negligible in the SM. This continues to hold in the MSSM for small $\tan\beta$. At large $\tan\beta$, depending on the Higgs sector there is a peculiar phenomenology already for minimal flavour violation, see e.g. [21,22,23,24,25,26,27,28,29]. For detailed studies of large-$\tan\beta$ effects beyond minimal flavour violation we refer to [30,31,32,33]. As large $\tan\beta \sim 30 - 60$ is separately reviewed at this conference [34], in the remainder we will largely restrict ourselves to small $\tan\beta$.

3 Phenomenology

We turn to a discussion of a number of flavour observables that can provide, at present or in the foreseeable future, potential for constraining or discovering supersymmetry. Schematically, a decay rate provides a constraint

$$
\begin{aligned}
&|\mathcal{A}_{\rm SUSY}|^2 + 2\,{\rm Re}\,\mathcal{A}^*_{\rm SUSY}\mathcal{A}_{\rm SM} \\
&\quad = \Gamma_{\rm exp}(1 \pm \Delta^{(\rm exp)}) - |\mathcal{A}_{\rm SM}|^2 (1 \pm \Delta^{(\rm SM)}).
\end{aligned}
$$

The right-hand side often involves a cancellation: the flavour observables measured so far are consistent with no SUSY contributions. Hence the goal is precision, reducing the errors $\Delta^{(\rm exp)}$ and $\Delta^{(\rm SM)}$ as much as possible, while for the left-hand side the rough dependence on SUSY parameters is enough. Other observables have similar expressions. The focus below is on hadronic observables.

3.1 Mixing

3.1.1 $K - \bar K$ mixing

$K - \bar K$ oscillations played a role in estimating the charm quark mass before its observation, as well as in the discovery of (indirect) CP violation, later giving information on the CP-violating phase in the CKM matrix. ΔM_K and ϵ_K also provide a classic constraint on supersymmetric flavour (see e.g. [35,36]). The mass difference ΔM_K and the CP-violating parameter ϵ_K follow from the effective $\Delta F = 2$ Hamiltonian,

$$\Delta M_K \propto 2 \sum_i B_i \,\mathrm{Re}\, C_i, \tag{28}$$

$$\epsilon_K \propto \frac{e^{i\pi/4}}{\sqrt{2}\Delta M_K} \sum_i B_i \,\mathrm{Im}\, C_i, \tag{29}$$

where $B_i = \langle K|Q_i|\bar{K}\rangle$. The hadronic matrix elements B_i contain low-energy QCD effects and require nonperturbative methods such as (numerical) lattice QCD, see e.g. [37,39,38]. Moreover, ΔM_K is afflicted by long-distance contributions which are believed to be subdominant but difficult to estimate. Nevertheless, in view of the strong CKM suppression of the SM contribution, even a rough estimate of the B_i translates into strong constraints on $s \to d$ flavour violation parameters, leading to bounds (assuming gluino-squark dominance and absence of cancellations)

$$|(\delta^d_{ds})_{LL}|, |(\delta^d_{ds})_{RR}| < \mathcal{O}(10^{-2}), \tag{30}$$

$$|(\delta^d_{ds})_{LR}|, |(\delta^d_{ds})_{RL}| < \mathcal{O}(10^{-3}), \tag{31}$$

$$|(\delta^d_{ds})_{LL} \cdot (\delta^d_{ds})_{RR}| < \mathcal{O}(10^{-7}), \tag{32}$$

for $M \sim m_{\tilde{q}} \sim m_{\tilde{g}} \sim 200$ GeV [15,40,41], a well-known aspect of the "SUSY flavour problem". The constraints become weaker as the SUSY scale is increased, scaling roughly like M, as is evident from (16)–(20). At any rate, this "problem" looks less severe when considering that the corresponding CKM factor $V^*_{td}V_{ts} = \mathcal{O}(10^{-4})$ is also much smaller than its 'generic' value $\mathcal{O}(1)$, and that the LR δ parameters are $\mathcal{O}(v/M)$. Indeed the problem is completely removed, for instance, in the framework of minimal flavour violation [42,43,44,45].

3.1.2 $B_d - \bar{B}_d$ and $B_s - \bar{B}_s$ mixing

Here the mixing amplitudes

$$\mathcal{A}(\bar{B}_q \to B_q) \propto M^q_{12} - \frac{i}{2}\Gamma^q_{12} \tag{33}$$

($q = d, s$) are completely short-distance dominated. Hence the theoretical expression

$$\Delta M_{B_q} \propto |M^q_{12}| \sim f^2_{B_q} |\sum B_i C_i|, \tag{34}$$

where f_{B_q} are decay constants, and B_i again parameterize hadronic matrix elements, can be directly compared to the experimental values [46,47]

$$\Delta M_{B_q} = (0.507 \pm 0.004)\,\mathrm{ps}^{-1}, \tag{35}$$

$$\Delta M_{B_s} = (17.77 \pm 0.10 \pm 0.07)\,\mathrm{ps}^{-1}. \tag{36}$$

In both cases, the theory error is fully dominated by f_{B_q}. For instance, $\Delta M^{\mathrm{SM}}_{B_s} \approx 16 \ldots 27\,\mathrm{ps}^{-1}$ seems realistic depending on the values of f_{B_s} [48,49,50]. This is consistent with the experimental value (which is however on the low side). Combined with the fact that the remaining (perturbative, non-CKM) uncertainties are at the 1–2 percent level, this underlines the importance of the ongoing efforts to obtain these nonperturbative parameters on the lattice with a high precision.

On the other hand, the weak phase $\phi_d = \arg M^d_{12}$ governs mixing-decay interference, hence can be extracted cleanly from the time-dependent CP asymmetry in $B \to J/\psi K_S$ decay (with theoretical uncertainties of $\mathcal{O}(1-2)$ %), giving [46]

$$\sin\phi_d = 0.675 \pm 0.026.$$

In the SM, $\phi_d = 2\beta$, but this does not hold in the presence of new flavour violation. Constraints analogous to (30)-(32) follow from $B_d - \bar{B}_d$ mixing (see e.g. [51]).

The mixing phase ϕ_s in the B_s system will be measured at LHCb from the analogous asymmetry in $B_s \to J/\psi\,\phi$. In the standard model $\phi_s \approx 0$, and any mixing-induced asymmetry in this mode would be a crystal clear signal of new physics.

The impact of $B_s - \bar{B}_s$ mixing data on the MSSM [33,52,53] is considered below in subsection 3.5, and in section 4 in the context of SUSY GUTs.

3.1.3 $D - \bar{D}$ mixing

The observation of $D-\bar{D}$ oscillations in 2007 at the B-factories [54,55,56] provides a constraint, which is unfortunately difficult to quantify because the SM contribution to M_{12} is completely long-distance dominated and rather uncertain. What is certain is that it has negligible weak phase, hence mixing-induced CP violation in D decays would signal non-SM physics. For the time being, an upper bound on $|M_{12}|$ can be obtained from data and used to put constraints on up-type δ parameters analogous to the cases discussed above [57,58,59,60].

3.2 Rare K decays

The decays $K^+ \to \pi^+\nu\bar{\nu}$ and $K_L \to \pi^0\nu\bar{\nu}$ are almost unique in that they are essentially free of hadronic uncertainties. In the SM context, the two modes provide a clean and independent determination of the unitarity triangle [61,62,63,64]— once they will have been measured precisely, hopefully, at CERN NA48/III and at JPARC (the SM branching fractions are $\mathcal{O}(10^{-11})$ and $\mathcal{O}(10^{-10})$, respectively). In general, they rather selectivly probe the FCNC vertices of the Z boson. In view of the particular, $SU(2)$-breaking structure of the leading Zsd vertex, this implies a specific sensitivity to certain combinations of LR and RL δ-parameters, even in the presence of general flavour violation [66,67,68,69,70,71]. In a perturbative expansion in δ's,

$$\mathcal{A}(K \to \pi\nu\bar{\nu})^{\mathrm{SUSY}} \propto (\delta^u_{dt})_{LR}(\delta^u_{ts})_{RL}, \tag{37}$$

where $(\delta^u_{dt(st)})_{LR} = (\delta^u_{td(ts)})^*_{RL}$ are related to up-squark (mass)^2 matrix elements in a certain non-super-CKM basis [67]. This is an example where the (generically)

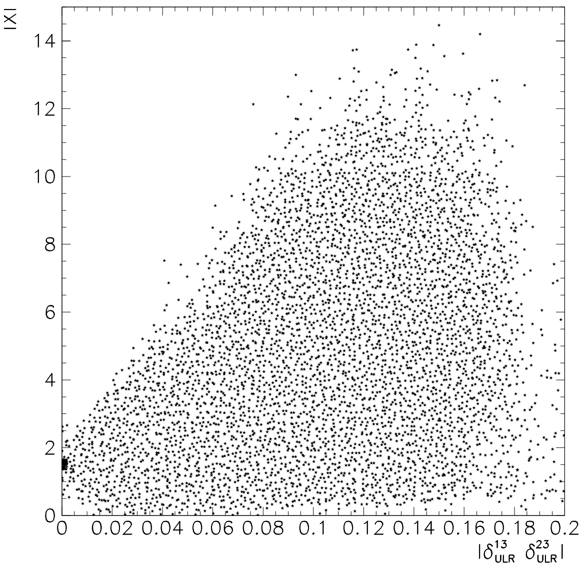

Fig. 1. SUSY contribution X to $\mathcal{A}(K \to \pi\nu\bar{\nu})$. Figure taken from [70]

leading effect arises at second order in the mass insertions. A systematic numerical analysis [70] (see also [71]) shows that this parametric dependence continues to hold beyond the perturbative expansion, and even in the presence of large contributions from box diagrams (Fig. 1). Moreover the SM hierarchy between the charged and neutral modes may be reversed, the latter enhanced by 1 to 2 orders of magnitude, and the bound following from isospin [65] saturated. Complementary probes of the Z-penguin amplitude are provided by the modes $K_L \to \pi^0 e^+e^-$ and $K_L \to \pi^0\mu^+\mu^-$, which still are theoretically quite clean [71, 72].

3.3 Leptonic B decays

Among B-decays, the modes $B^+ \to \ell^+\nu$ and $B_d, B_s \to \ell^+\ell^-$ are the theoretically cleanest. The former proceeds through a W^+ tree-level diagram in the SM. Significant corrections may occur due to charged-higgs boson exchange, which is present in the MSSM [73, 74, 27] but becomes relevant only at large values of $\tan\beta$. In this case, the latter mode can be enhanced by an order of magnitude or more, which provides a serious constraint on certain large-$\tan\beta$ scenarios. (The branching fraction scales with the sixth power of $\tan\beta$.) At small $\tan\beta$, it will receive more moderate contributions via the Z-penguin and the operator Q_{10A}. In the SM one has [75, 76]

$$BR(B_s \to \mu^+\mu^-) = (3.51 \pm 0.50) \times 10^{-9},$$

where the bulk of the hadronic uncertainties has been eliminated by normalizing to ΔM_s. (The error will become even smaller with improved lattice predictions for $\hat{B}_{B_s}$.) In spite of this mode being so rare, LHCb, ATLAS, and CMS expect to collect a combined few hundred SM events after five years or so of running. Hence this mode will play an important role also at small $\tan\beta$.

3.4 Semileptonic and radiative decays

3.4.1 Inclusive $B \to X_s\gamma$

While being loop-induced, the inclusive decay $B \to X_s\gamma$ has a relatively large branching ratio and is well measured, with a world average [77] of

$$BR(\bar{B} \to X_s\gamma)_{\rm exp} = (3.55 \pm 0.24^{+0.09}_{-0.10} \pm 0.03) \times 10^{-4}$$

(with a 1.6 GeV lower cut on the photon energy). In the SM, the decay amplitude receives its dominant contributions from $W - t$ loops entering through the magnetic operator $Q_{7\gamma}$ and $W - c$ loops entering through loop contractions of the tree operators $Q_{1,2}$. Both contributions are of comparable size and opposite in sign. Other operators are subdominant. The corresponding state-of-the-art NNLO-QCD prediction reads [78] (see [77] for a review and discussion of uncertainties)

$$BR(\bar{B} \to X_s\gamma)_{\rm SM} = (3.15 \pm 0.23) \times 10^{-4},$$

slightly more than 1σ below the experiment. Supersymmetric effects have been investigated thoroughly in the literature both in minimal flavour violation [79, 80, 81, 82, 83, 84, 25, 28, 45] and beyond [85, 86, 87]. They enter chiefly through $Q_{7\gamma}$, resulting in a large sensitivity to the parameter $(\delta^d_{sb})_{LR}$. (Contributions via $Q'_{7\gamma}$ do not interfere with the SM contribution in an inclusive process due to the opposite chirality of the produced s-quark, hence generally have a small impact.) The full one-loop SUSY contribution may involve a cancellation between charged-Higgs-top loops, which are always of the same sign as the SM piece, and squark-higgsino as well as squark-gaugino loops, which (in general) carry an arbitrary complex phase.

3.4.2 Inclusive $B \to X_s\ell^+\ell^-$

Their sensitivity to semileptonic operators like Q_{9V} and Q_{10A} makes the rare $b \to s\ell^+\ell^-$ transitions a complementary and more complex test of the underlying theory than the radiative ones. The three-body decays allow to study non-trivial observables such as the dependence on the kinematics of the decay products. In the absence of large statistics, partially integrated spectra such as the dilepton mass spectrum or the angular distribution can be explored that are amenable to a clean theoretical description for a dilepton invariant mass below the charm resonances.

In the minimally flavour-violating MSSM the Wilson coefficients C_{9V} and C_{10A} are only slightly affected and corrections to the decay distributions do not exceed the 30% level [88, 89]. At large $\tan\beta$, additional contributions to $b \to s\mu^+\mu^-$ arise from the chirality-flipping operators $(\bar{s}_L b_R)(\bar{\mu}_L\mu_R)$ and $(\bar{s}_L b_R)(\bar{\mu}_R\mu_L)$ that are suppressed by powers of the muon mass but enhanced by $(\tan\beta)^3$. In practice, these contributions are however bounded from above [90, 91, 92] by the experimental constraints on $B_s \to \mu^+\mu^-$ and turn out

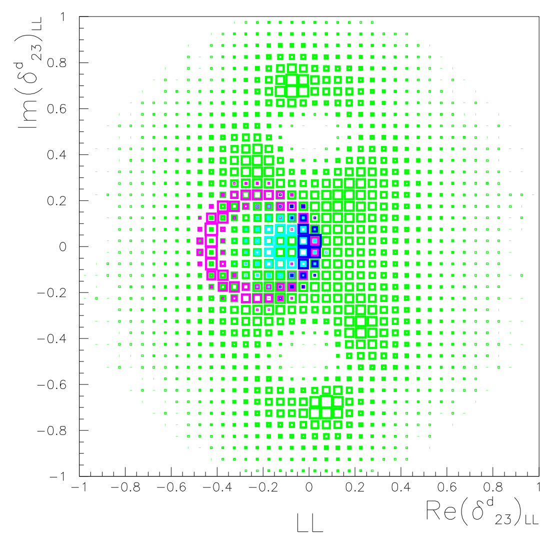

Fig. 2. Combined constraints on $(\delta^d_{sb})_{LL}$. The allowed regions due to $B \to X_s\gamma$, $B \to X_s\ell^+\ell^-$, and $B_s - \bar{B}_s$ mixing are shown in violet, light blue, and green, respectively; the combined constraint is in dark blue. A common sparticle mass of 350 GeV and $\mu = -350$ GeV is assumed. Figure taken from [94]

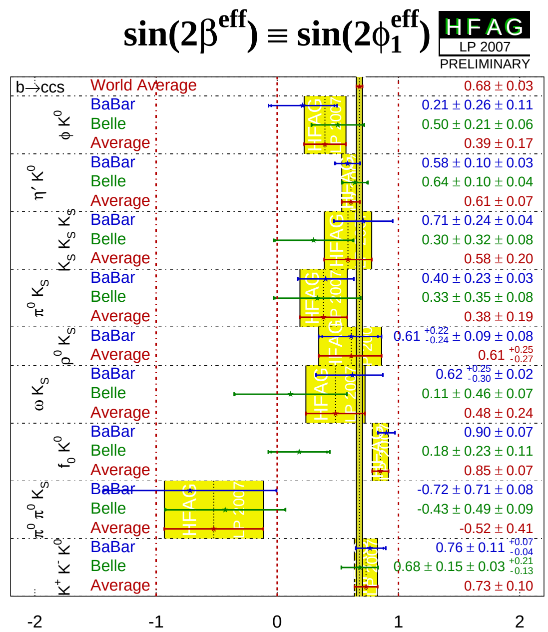

Fig. 3. Mixing-induced CP violation in charmless $b \to s$ transitions. Figure taken from [46]

to be subleading. Merging the information on $B \to X_s\ell^+\ell^-$ with the one on $B \to X_s\gamma$, one can thus infer that the sign of the $b \to s\gamma$ amplitude should be SM-like [93]. In the general MSSM a simultaneous use of the $B \to X_s\ell^+\ell^-$ and $B \to X_s\gamma$ constraints leads to stringent limits on $(\delta^d_{sb})_{LL}$ and $(\delta^d_{sb})_{LR}$ in the complex plane [87,94].

3.5 Combined constraints

The constraints on the flavour-violating parameters become more powerful when the interplay of several observables with different parametric sensitivities is considered. (This was done in many of the works referred to above and below.) For $b \to s$ transitions, Fig. 2 (taken from [94]) shows how the measurements of from $B \to X_s\gamma$, $B \to X_s\ell^+\ell^-$, and ΔM_s coact to constrain the parameter $(\delta^d_{sb})_{LL}$, leaving a much smaller allowed region than each individual observable.

3.6 Charmless hadronic decays

Two-body exclusive nonleptonic decays $B \to M_1M_2$ are sensitive to all of the operators $Q^{(\prime)}_{1\ldots10,7\gamma,8g}$, while offering a large number $\mathcal{O}(100)$ of observables, including many CP-violating ones. They (and exclusive modes in general) also become increasingly attractive on the grounds of their accessibility at hadron machines such as the LHC. On the theoretical side, the factor limiting the precision are the hadronic matrix elements $\langle M_1M_2|Q_i|B\rangle$, which involve nonperturbative QCD in a way that is presently not surmountable in lattice QCD. Systematic methods are, however, available, based on expansions about the limit of $SU(3)$ flavour symmetry or about the heavy-b-quark limit. m_s/Λ and Λ/m_b, respectively, are the expansion parameters. In fact, the (QCD) factorization formulae [95,96,97] for the hadronic matrix elements that follow from the heavy-quark limit respect the $SU(3)$ flavour symmetry up to well-defined (so-called "factorizable") corrections at the leading power and perturbative order, and perturbative QCD corrections do not alter this picture much. Higher orders in Λ/m_b are generally not under control, with certain (important) exceptions.

3.6.1 Time-dependent CP asymmetries

One class of observables that has received much interest in recent years are time-dependent CP asymmetries in decays into CP eigenstates,

$$\frac{BR(\bar{B}^0(t) \to f) - BR(B^0(t) \to f)}{BR(\bar{B}^0(t) \to f) - BR(B^0(t) \to f)} \equiv S_f \sin(\Delta m_B t) - C_f \cos(\Delta m_B t). \quad (38)$$

In the case of charmless $b \to s$ transitions, within the SM one expects $\eta_f S_f \approx \phi_d = 2\beta$, based on the dominance of the QCD penguin amplitude, which has vanishing weak phase. (η_f denotes the CP quantum number of the final state. Neither equality holds in the MSSM (beyond minimal flavour violation). Fig. 3

shows the data for a number of modes. It is conspicuous that in general, the $\eta_f S_f \equiv \sin(2\beta^{\rm eff})(f)$ lies below the value of $\sin 2\beta$. In fact, in the SM a QCD factorization calculation [98] of the corrections for the two-body final states ϕK^0, $\eta' K^0$, $\pi^0 K_S$, $\rho^0 K_S$, ωK_S due to neglected operators $Q_{1,2}$ shows a (small) positive shift in all cases except $\rho^0 K_S$. Supersymmetric contributions to $b \to s$ penguin transitions are well studied, see for instance [99,100,101,102,103,104,105, 106,107,108,109,87,110,111,112]. Often the most important operator is Q_{8g}, or even Q'_{8g}, which will interfere with the SM contributions in exclusive processes. An interesting possibility is to attribute the pattern of the deviations to constructive and destructive interference between operators of different quark chiralities, depending on the parities of the final-state particles [103]. We stress that while the present data does not appear to show any significant deviations from SM expectations, the situation in Fig. 3 clearly illustrates the discovery potential of these hadronic observables.

3.7 Lepton flavour violation

We do not have time (space) to cover lepton flavour violation in detail, which is fortunately covered in many other talks at this conference. In the SM, even when introducing the minimal dimension-5 operator to allow for neutrino masses and mixings, lepton-flavour violating processes such as $\tau \to \mu\gamma$ are rendered unobservably rare by the tiny neutrino mass splittings and the corresponding near-perfect GIM cancellation. The situation is very different in the MSSM because of the presence of lepton flavour violation at the renormalizable level, in the sneutrino and charged slepton mass matrices. Theoretically, the SUSY effects in $\ell_i \to \ell_j\gamma$ can be captured in operators analogous to $Q_{7\gamma}$ figuring in the discussion of $B \to X_s\gamma$ above. The experimental upper bounds can likewise be converted into knowledge about the slepton mass matrices.

4 Probing the GUT scale

Concrete assumptions about the SUSY-breaking mechanism (gravity mediation, gauge mediation, etc.) and possible UV completion (such as a SUSY grand-unified theory, minimal flavour violation, etc.) may imply patterns in $\mathcal{L}_{\rm soft}$ relating different δ parameters that can be tested against the general constraints applying to them, or can be further used in making specific predictions for low-energy observables and their correlations. Recent trends in SUSY models are reviewed in other talks at this conference [114,115]. One of the most intriguing aspects of SUSY grand unified theories (GUTs), which also demonstrates the power of flavour observables to probe even superhigh scales, is the possibility of relations between hadronic and leptonic flavour violation (see e.g. [117,118,119,120,121, 122,123,124,125,126,127,128,129,130,131,132,133,134]). The effect we are considering here is the following [8, 117,118]. Assume that SUSY breaking is effected at a scale beyond the GUT scale, for instance at the Planck scale, and that it is flavour-blind, at least approximately, such that one has a universal scalar mass parameter m_0^2 and a universal A-parameter a_0. For definiteness, assume simple $SO(10)$ unification such that there is only one sfermion multiplet for each generation. Radiative corrections due to the unified gauge coupling will correct the masses of the three 16's of sfermions in the same way, while the large top Yukawa coupling will selectively suppress the masses of one multiplet,

$$m^2_{16_{1,2}} = m_0^2 + \epsilon, \qquad m^2_{16_3} = m_0^2 + \epsilon - \Delta, \qquad (39)$$

where $\epsilon \propto g^2$ and $\Delta \propto y_t^2 m_0^2,\, y_t^2 a_0^2$. Eq. (39) holds in a basis where the up-type Yukawa matrix is diagonal. (Further contributions will be present if there are additional large Yukawa couplings, such as for large $\tan\beta$. The expressions then become more complicated.) In this fashion, the large top Yukawa coupling affects also the right-handed down-type squark masses and the slepton masses, which is very different from the situation in the MSSM (or below the GUT scale). Now the relevant sfermion basis for low-energy physics (the super-CKM basis) is the one where the down-type and leptonic Yukawas are diagonal. In such a basis the members of the 16_3 selected by y_t will consist of mixtures of the superpartners of fermions of different flavours, and FCNC SUSY vertices appear. Moreover, these vertices will be correlated between the hadronic and leptonic sectors. This kind of mechanism gained particular attraction after the observation of large leptonic mixing in neutrino oscillations [124]. Assuming a minimal-$SU(5)$-type embedding of the MSSM into $SO(10)$, and under certain assumptions about the generation of seesaw neutrino masses, one has

$$Y_D = U^T_{\rm PMNS}\, {\rm diag}(y_d, y_s, y_b)\, V^\dagger_{\rm CKM}. \qquad (40)$$

This expression should be most robust in the $(2,3)$ sub-block. The atmospheric neutrino mixing angle then appears in the gluino-squark couplings

$$\mathcal{L}_{\rm soft} \supset U^{ij}_{\rm PMNS}\, \tilde d^*_{Rj}\, \tilde g T^A\, d_{Ri}, \qquad (41)$$

yielding potentially spectacular effects in observables like ΔM_{B_s} [108,129,130,131]. The model is parameterized by four parameters m_0, a_0, $m_{\tilde g}$, and μ, and a value of $\tan\beta$ around $2-3$ to maintain perturbativity. Fig. 4 compares $BR(\tau \to \mu\gamma)$ with the correction to ΔM_{B_s}. It is evident that the measurement of ΔM_{B_s} provides a nontrivial and quantifiable constraint on the lepton-flavour-violating mode, providing (one of many) illustration(s) of the possibility to probe very fundamental scales with flavour-violating observables, beyond what is possible with knowing the particle spectrum alone.

5 Conclusion

We have reviewed some of the possible indirect signals of SUSY in flavour-physics and the most important

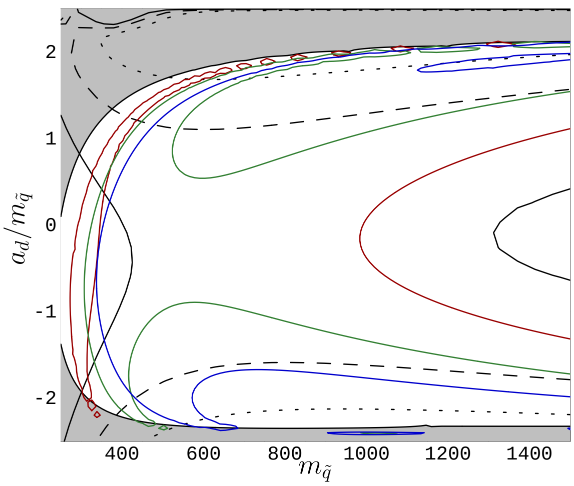

Fig. 4. Contours of constant SUSY contribution to ΔM_{B_s}, in units of the SM value, and of constant $BR(\tau \to \mu\gamma)$, in a slice of parameter space. Here a_d has a simple relation to a_0, $m_{\tilde{g}} = 250$ GeV, $\tan\beta = 3$. The solid black, dashed, and dotted contours correspond to $|\Delta M_{B_s}^{(\mathrm{NP})}|/\Delta M_{B_s}^{(\mathrm{SM})} =$ 0.5, 2, 5, respectively. The red, green, and blue contours correspond to the experimental upper bound on $BR(\tau \to \mu\gamma)$ for $\mu = -300$, $\mu = -450$, $\mu = -600$ GeV, respectively. (ΔM_{B_s}) is independent of μ in the approximation used. Figure from [131]

constraints that current data imposes. We have also discussed some modes and signals that will be measured in the near future and are likely to be affected by TeV-scale supersymmetry. Neither set is complete, for instance we did not have room to discuss exclusive semileptonic B-decays, which are under active theoretical and experimental investigation. To close on a positive note, we can hope that soon the era of putting bounds and constraints will give way to a phase of actual measurements of MSSM parameters. The interplay between direct and indirect observables should be useful, as it was in the construction of the SM. For the longer term, it is then conceivable that after such a phase indirect probes will take the center stage again. While it may be the past experience that "there is interesting physics at all scales", the peculiar gentleness of SUSY quantum corrections may mean that this cannot be extrapolated further. In SUSY, the GUT or Planck scales can actually be "close" as far as indirect observables are considered (as argued in the previous section), even if being at great distance from the point of view of direct detection.

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