B-decay CP-asymmetries in SUSY with a $U(2)^3$ flavour symmetry

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We study CP asymmetries in rare B decays within supersymmetry with a $U(2)^3$ flavour symmetry, motivated by the SUSY flavour and CP problems, the hierarchies in the Yukawa couplings and the absence so far of any direct evidence for SUSY. Even in the absence of flavour-blind phases, we find potentially sizable CP violating contributions to $b \to s$ decay amplitudes. The effects in the mixing-induced CP asymmetries in $B \to \phi K_S$ and $B \to \eta' K_S$, angular CP asymmetries in $B \to K^* \mu^+ \mu^-$ and the direct CP asymmetry in $B \to X_s \gamma$ can be in the region to be probed by LHCb and next generation B factories. At the same time, these effects in B decays are compatible with CP violating contributions to meson mixing, including a non-standard B_s mixing phase hinted by current tensions in the CKM fit mostly between $S_{\psi K_S}, \epsilon_K$ and $\Delta M_{B_s}/\Delta M_{B_d}$.

1. Introduction

Weak scale supersymmetry has long been the leading candidate for physics beyond the Standard Model (SM), but it is increasingly challenged by the lack of any clear experimental evidence in its favour. Arguably, the most important challenges to be explained are the success of the CKM description of flavour violation, the tight bounds on flavour-blind CP violation from electric dipole moment (EDM) searches and the tightening sparticle mass bounds from LHC.

A promising approach to address these three problems is to make use of the fact that the third generation of quarks is special, since it has sizable Yukawa couplings and is weakly mixed with the first two generations. If the first two generations of quarks and squarks form doublets under a $U(2)$ flavour symmetry commuting with the gauge group, the flavour problem is ameliorated while at the same time the hierarchies in the Yukawa couplings can be (at least partially) understood [1, 2]. Furthermore, a flavour symmetry for the light quark generations fits well to a hierarchy in the spectrum of squarks, with the third generation partners being light and the first two generation ones in the multi-TeV region [3, 4, 5, 6]. This has two virtues, in addition to further ameliorating the flavour problem: First, it solves the CP problem, decoupling the contributions of flavour-blind phases to the observable EDMs, which involve first generation fermions [7]. Second, it evades the strongest sparticle direct search bounds. While the lower bounds on first two generation squark masses from LHC are approaching a TeV, third generation squarks can still be in the few-hundred GeV region for gluino masses above about 600 GeV, depending on the gluino decay branching ratios and on the lightest neutralino mass.[1]

Motivated by these facts, a $U(2)^3$ symmetry has been considered in [10] together with a hierarchical squark spectrum and has been shown to solve, in addition to the above-mentioned problems, tensions present in the fits of the CKM matrix by contributing to CP violation in meson mixing. In this way, a non-standard B_s mixing phase was predicted, as currently hinted by Tevatron data [11, 12].

The purpose of this article is to extend the analysis in [10] to $\Delta F = 1$ processes, i.e. rare decays, studying in particular possible signatures of CP violation correlated with the predicted CP violation in meson mixing ($\Delta F = 2$). Contrary to the $\Delta F = 2$ sector, where the pattern of deviations from the SM identified in [10] is unambiguously dictated by the $U(2)^3$ symmetry, the predictions of $\Delta F = 1$ observables are more model dependent. In this analysis we concentrate on a framework with moderate values of $\tan\beta$ and, in order to establish a link between $\Delta F = 2$ and $\Delta F = 1$ CP-violating observables, we take small flavour-blind phases and assume the dominance of gluino-mediated flavour-changing amplitudes.

2. Setup

We briefly review the $U(2)^3$ framework, referring to [10] for details and derivations. We consider a global flavour symmetry

$$G_F = U(2)_Q \times U(2)_u \times U(2)_d \tag{1}$$

broken *minimally* by three spurions, transforming as $\Delta Y_u = (2, \bar{2}, 1)$, $\Delta Y_d = (2, 1, \bar{2})$ and $V = (2, 1, 1)$, respectively. This breaking pattern leads to non-minimal flavour violation in the down quark-squark-gluino vertex,

$$(\bar{d}_{L,R} W^d_{L,R} \tilde{d}_{L,R})\tilde{g}. \tag{2}$$

which can be approximately written as

$$W^d_L = \begin{pmatrix} c_d & \kappa^* & -\kappa^* s_L e^{i\gamma_L} \\ -\kappa & c_d & -c_d s_L e^{i\gamma_L} \\ 0 & s_L e^{-i\gamma_L} & 1 \end{pmatrix}, \qquad W^d_R = \mathbb{1}\,, \tag{3}$$

where $\kappa = c_d V_{td}/V_{ts}$ and c_d is fixed by the CKM matrix, which in these conventions reads

$$V_{\rm CKM} = \begin{pmatrix} 1-\lambda^2/2 & \lambda & s_u s e^{-i\delta} \\ -\lambda & 1-\lambda^2/2 & c_u s \\ -s_d s\, e^{i(\phi+\delta)} & -s c_d & 1 \end{pmatrix}, \tag{4}$$

[1]For a first LHC analysis with about 1 fb^{-1} of integrated luminosity, see [8, 9].

where $s_u c_d - c_u s_d e^{-i\phi} = \lambda e^{i\delta}$. A direct fit of this parametrization to tree-level observables results in

$$s_u = 0.086 \pm 0.003\,, \qquad s_d = -0.22 \pm 0.01\,, \tag{5}$$
$$s = 0.0411 \pm 0.0005\,, \qquad \phi = (-97 \pm 9)^\circ\,. \tag{6}$$

The mixing s_L and the phase[2] γ_L in (3), on the other hand, are free parameters. Defining

$$\xi_L = \frac{c_d s_L}{|V_{ts}|} e^{i\gamma_L}\,, \tag{7}$$

the mixing amplitudes in the K, B_d and B_s meson systems are modified as follows,

$$M_{12}^K = \left(M_{12}^K\right)_{\rm SM}^{(tt)} \left(1 + |\xi_L|^4 F_0\right) + \left(M_{12}^K\right)_{\rm SM}^{(tc+cc)}\,, \tag{8}$$
$$M_{12}^{B_{d,s}} = \left(M_{12}^{B_{d,s}}\right)_{\rm SM} \left(1 + \xi_L^2 F_0\right)\,, \tag{9}$$

where $F_0 > 0$ is a loop function, predicting a universal shift in the B_d and B_s mixing phases. The mixing induced CP asymmetries in $B \to \psi K_S$ and $B_s \to \psi\phi$ are thus

$$S_{\psi K_S} = \sin\left(2\beta + \phi_\Delta\right)\,, \qquad S_{\psi\phi} = \sin\left(2|\beta_s| - \phi_\Delta\right)\,, \tag{10}$$

where $\phi_\Delta = \arg\left(1 + \xi_L^2 F_0\right)$.

The current tensions in the CKM fit arising mostly between $S_{\psi K_S}$, $|\epsilon_K|$ and $\Delta M_s/\Delta M_d$ can be accommodated by the shifts in $S_{\psi K_S}$ and $|\epsilon_K|$. Assuming this to be the case, a global fit of the CKM matrix to $\Delta F = 2$ observables including the supersymmetric contributions leads to the following predictions at the 90% C.L.,

$$0.8 < |\xi_L| < 2.1\,, \tag{11}$$
$$-9^\circ < \phi_\Delta < -1^\circ\,, \tag{12}$$
$$-86^\circ < \gamma_L < -25^\circ \quad \text{or} \quad 94^\circ < \gamma_L < 155^\circ\,, \tag{13}$$
$$0.05 < S_{\psi\phi} < 0.20,\,. \tag{14}$$

with gluino and left-handed sbottom masses necessarily below $1\div1.5$ TeV. The ambiguity in γ_L stems from the fact that the $\Delta F = 2$ observables are only sensitive to the phase of ξ_L^2, and thus to $2\gamma_L$.

The presence of the phase γ_L leads to contributions to CP asymmetries in B physics. Additional contributions can arise from flavour-blind phases, such as the phase of the μ term or the trilinear couplings. With hierarchical sfermions, such phases are not required to be small to meet the EDM constraints [7]. Still, to concentrate on the genuine $U(2)^3$ effects, we will take flavour blind phases to be absent in the following and thus focus on B physics.

[2] Our phase γ_L is equivalent to the phase γ in [10].

3. $\Delta B = 1$ effective Hamiltonian

The part of the $b \to s$ effective Hamiltonian sensitive to NP in our setup reads

$$\mathcal{H}_{\text{eff}} = \frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb} \sum_{i=3}^{10} C_i O_i + \text{ h.c.}\,, \tag{15}$$

$$\begin{aligned}
O_3 &= (\bar{s} P_L b) \textstyle\sum_q (\bar{q} P_L q) & O_4 &= (\bar{s}_\alpha P_L b_\beta) \textstyle\sum_q (\bar{q}_\beta P_L q_\alpha) \\
O_5 &= (\bar{s} P_L b) \textstyle\sum_q (\bar{q} P_R q) & O_6 &= (\bar{s}_\alpha P_L b_\beta) \textstyle\sum_q (\bar{q}_\beta P_R q_\alpha) \\
O_7 &= \frac{e}{16\pi^2} m_b (\bar{s} \sigma_{\mu\nu} P_R b) F_{\mu\nu} & O_8 &= \frac{g_s}{16\pi^2} m_b (\bar{s} \sigma_{\mu\nu} P_R b) G_{\mu\nu} \\
O_9 &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_L b)(\bar{l} \gamma_\mu l) & O_{10} &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_L b)(\bar{l} \gamma_\mu \gamma_5 l)
\end{aligned}$$

The QCD penguin operators $O_{3\ldots6}$ are relevant for the CP asymmetries in $b \to s\bar{s}s$ penguin decays to be discussed below. In the case of hierarchical sfermions, the box contributions are mass-suppressed and we only have to consider gluon penguins, which contribute in a universal way as

$$C_3 = C_5 = -\tfrac{1}{3} C_4 = -\tfrac{1}{3} C_6 \equiv C_G\,. \tag{16}$$

The analogous photon penguin can be neglected in non-leptonic decays since it is suppressed by α_{em}/α_s, but it can contribute to C_9. There is no effect in C_{10}, which remains SM-like.

In the MSSM, the $\Delta F = 1$ effective Hamiltonian receives contributions from loops involving charginos, neutralinos, charged Higgs bosons or gluinos. In the following, we will concentrate for simplicity on gluino contributions, which are always proportional to the complex ξ_L in $U(2)^3$. Among the remaining contributions, some are proportional to ξ_L while some are real in the absence of flavour blind phases. Their omission does not qualitatively change our predictions for CP asymmetries, but we stress that the real contributions can have an impact in particular on the branching ratios to be considered below.

The gluino contributions to C_G and C_9 are

$$C_G = -\xi_L \, \frac{\alpha_s}{\alpha_2} \, \frac{\alpha_s}{4\pi} \, \frac{m_W^2}{m_{\tilde{b}}^2} \, \frac{13}{108} f_G(x_{\tilde{g}})\,, \qquad C_9 = \xi_L \, \frac{\alpha_s}{\alpha_2} \, \frac{m_W^2}{m_{\tilde{b}}^2} \, \frac{2}{27} f_\gamma(x_{\tilde{g}})\,, \tag{17}$$

where here and throughout, $x_{\tilde{g}} = m_{\tilde{g}}^2/m_{\tilde{b}}^2$ and all the loop functions are defined such that $f_i(1) = 1$ with the exact form given in appendix A.

The main difference concerning the magnetic and chromomagnetic Wilson coefficients C_7 and C_8 is that here $\tilde{b}_L$-$\tilde{b}_R$ mass insertion contributions are important, while for the other Wilson coefficients they were chirality-suppressed. The gluino contributions read

$$C_7 = -\xi_L \, \frac{\alpha_s}{\alpha_2} \, \frac{m_W^2}{m_{\tilde{b}}^2} \, \frac{1}{27} \left[f_7(x_{\tilde{g}}) + 2 \frac{\mu \tan\beta - A_b}{m_{\tilde{g}}} g_7(x_{\tilde{g}}) \right]\,, \tag{18}$$

$$C_8 = -\xi_L \, \frac{\alpha_s}{\alpha_2} \, \frac{m_W^2}{m_{\tilde{b}}^2} \, \frac{5}{36} \left[f_8(x_{\tilde{g}}) + 2 \frac{\mu \tan\beta - A_b}{m_{\tilde{g}}} g_8(x_{\tilde{g}}) \right]\,. \tag{19}$$

Observable	SM prediction	Experiment	Future sensitivity
$\mathrm{BR}(B \to X_s\gamma)$	$(3.15 \pm 0.23) \times 10^{-4}$ [14]	$(3.52 \pm 0.25) \times 10^{-4}$	$\pm 0.15 \times 10^{-4}$
$A_{\mathrm{CP}}(b \to s\gamma)$	$(-0.6 \div 2.8)\%$ [15]	$(-1.2 \pm 2.8)\%$	$\pm 0.5\%$
$\mathrm{BR}(B \to X_d\gamma)$	$(1.54^{+0.26}_{-0.31}) \times 10^{-5}$ [16]	$(1.41 \pm 0.49) \times 10^{-5}$	
$S_{\phi K_S}$	0.68 ± 0.04 [17, 18]	$0.56^{+0.16}_{-0.18}$	± 0.02
$S_{\eta' K_S}$	0.66 ± 0.03 [17, 18]	0.59 ± 0.07	± 0.01
$\langle A_7 \rangle$	$(3.4 \pm 0.5) \times 10^{-3}$ [19]	–	
$\langle A_8 \rangle$	$(-2.6 \pm 0.4) \times 10^{-3}$ [19]	–	

Table 1: SM predictions, current experimental world averages [20] and experimental sensitivity at planned experiments [21, 22] for the B physics observables. $< A_{7,8} >$ are suitable angular CPV asymmetries in $B \to K^* \mu^+ \mu^-$.

For $m_{\tilde{g}} = m_{\tilde{b}} \equiv \tilde{m}$, we thus find at the scale m_W

$$C_G = -1.1 \times 10^{-4} \left(\frac{500\,\mathrm{GeV}}{\tilde{m}} \right)^2 \xi_L \,, \tag{20}$$

$$C_\gamma = 9 \times 10^{-3} \left(\frac{500\,\mathrm{GeV}}{\tilde{m}} \right)^2 \xi_L \,, \tag{21}$$

$$C_7 = -3.4 \times 10^{-3} \left(\frac{500\,\mathrm{GeV}}{\tilde{m}} \right)^2 \xi_L \left[1 + 2 \frac{\mu \tan\beta - A_b}{\tilde{m}} \right] , \tag{22}$$

$$C_8 = -1.3 \times 10^{-2} \left(\frac{500\,\mathrm{GeV}}{\tilde{m}} \right)^2 \xi_L \left[1 + 2 \frac{\mu \tan\beta - A_b}{\tilde{m}} \right] . \tag{23}$$

A model-independent consequence of the $U(2)^3$ symmetry is that the modification of $b \to s$ and $b \to d$ $\Delta F = 1$ amplitudes is *universal*, i.e. only distinguished by the same CKM factors as in the SM (exactly as in the $U(3)^3$, or MFV case [13]). Therefore, all the expressions for the $b \to s$ Wilson coefficients derived in this section are also valid for $b \to d$ processes.

4. B physics observables

4.1. $\mathrm{BR}(B \to X_q\gamma)$

The branching ratio of $B \to X_s\gamma$ is one of the most important flavour constraints in the MSSM in view of the good agreement between theory and experiment. Experimentally, the quantities

$$R_{bq\gamma} = \frac{\mathrm{BR}(B \to X_q\gamma)}{\mathrm{BR}(B \to X_q\gamma)_{\mathrm{SM}}} \tag{24}$$

are constrained to be

$$R_{bs\gamma} = 1.13 \pm 0.11\,, \qquad R_{bd\gamma} = 0.92 \pm 0.40\,, \tag{25}$$

using the numbers in table 1. In $U(2)^3$, one has $R_{bs\gamma} \approx R_{bd\gamma}$ just as in MFV, so the $b \to s\gamma$ constraint is more important and we will concentrate on it in the following.

Beyond the SM (but in the absence of right-handed currents), the branching ratio can be written as

$$\mathrm{BR}(B \to X_s\gamma) = R\left[\left|C_7^{\mathrm{SM,eff}} + C_7^{\mathrm{NP,eff}}\right|^2 + N(E_\gamma)\right], \tag{26}$$

where $R = 2.47 \times 10^{-3}$ and $N(E_\gamma) = (3.6 \pm 0.6) \times 10^{-3}$ for a photon energy cut-off $E_\gamma = 1.6$ GeV [23].

Considering only gluino contributions and setting $m_{\tilde{g}} = m_{\tilde{b}} \equiv \tilde{m}$, we find

$$R_{bs\gamma} = 1 \;+\; 2.2 \times 10^{-2}\left(\frac{500\,\mathrm{GeV}}{\tilde{m}}\right)^2 |\xi_L| \cos\gamma_L \left(1 + 2\frac{\mu\tan\beta - A_b}{\tilde{m}}\right). \tag{27}$$

As stressed in section 3, there are additional real contributions to the Wilson coefficient $C_{7,8}$ that can modify the branching ratio. In particular, there is a $\tan\beta$ enhanced chargino contribution proportional to the stop trilinear coupling, which can interfere constructively or destructively with the SM. Thus, with a certain degree of fine-tuning, the constraints in 25 can always be fulfilled. In our numerical analysis, we will require the branching ratio including *only SM and gluino contributions* to be within 3σ of the experimental measurement.

4.2. $\boldsymbol{A_{\mathrm{CP}}(B \to X_s\gamma)}$

The direct CP asymmetry in $B \to X_s\gamma$

$$A_{\mathrm{CP}}(B \to X_s\gamma) = \frac{\Gamma(\bar{B} \to X_s\gamma) - \Gamma(B \to X_{\bar{s}}\gamma)}{\Gamma(\bar{B} \to X_s\gamma) + \Gamma(B \to X_{\bar{s}}\gamma)}\,, \tag{28}$$

already constrained by the B factories as shown in table 1, will be measured by next generation experiments to a precision of 0.5%. On the theory side, the recent inclusion of "resolved photon" contributions reduced the attainable sensitivity to NP in view of the large non-perturbative SM contribution leading to [15]

$$-\,0.6\% < A_{\mathrm{CP}}(B \to X_s\gamma)_{\mathrm{SM}} < 2.8\% \tag{29}$$

compared to an earlier estimate of the short-distance part [24]

$$A_{\mathrm{CP}}(B \to X_s\gamma)_{\mathrm{SM}}^{\mathrm{SD}} = (0.44^{+0.24}_{-0.14})\%\,. \tag{30}$$

The new physics contribution can be written as

$$A_{\mathrm{CP}}(B \to X_s\gamma)_{\mathrm{NP}} \times R_{bs\gamma} = -0.29\mathrm{Im}(C_7^{\mathrm{NP}}) + 0.30\mathrm{Im}(C_8^{\mathrm{NP}}) - 0.99\mathrm{Im}(C_7^{\mathrm{NP}*}C_8^{\mathrm{NP}}), \tag{31}$$

valid for $E_\gamma = 1.85\,\mathrm{GeV}$.

Setting $m_{\tilde g} = m_{\tilde b} \equiv \tilde m$ and $R_{bs\gamma} = 1$, the gluino contributions to the CP asymmetry are

$$A_{\rm CP}(B \to X_s\gamma)_{\rm NP} = -1.74 \times 10^{-3} \left(\frac{500\,\mathrm{GeV}}{\tilde m}\right)^2 |\xi_L| \sin\gamma_L \left(1 + 2\frac{\mu\tan\beta - A_b}{\tilde m}\right) \quad (32)$$

we note that the CP asymmetry can have either sign due to the two solutions for γ_L allowed by $\Delta F = 2$, see (13).

4.3. $B \to K^*\mu^+\mu^-$

Angular CP asymmetries in $B \to K^*\mu^+\mu^-$ are sensitive probes of non-standard CP violation and will be measured soon at the LHCb experiment. In our framework, where right-handed currents are absent, the relevant observables are the T-odd CP asymmetries A_7 and A_8 as defined in [19].

For these observables, integrated in the low dilepton invariant mass region, we obtain the simple dependence on the Wilson coefficients, as usual to be evaluated at the scale m_b,

$$\langle A_7\rangle \times R_{\rm BR} \approx -0.71\,\mathrm{Im}(C_7^{\rm NP})\,, \quad (33)$$

$$\langle A_8\rangle \times R_{\rm BR} \approx 0.40\,\mathrm{Im}(C_7^{\rm NP}) + 0.03\,\mathrm{Im}(C_9^{\rm NP})\,, \quad (34)$$

where $R_{\rm BR}$ is the ratio between the full result for the CP-averaged branching ratio and the SM one [7], $R_{\rm BR} \approx 1$ in our framework. Although $\mathrm{Im}(C_7^{\rm NP})$ and $\mathrm{Im}(C_9^{\rm NP})$ can be of the same order, the contribution from $C_9^{\rm NP}$ is numerically suppressed and one will thus still have approximately

$$\langle A_8\rangle \simeq -0.56\langle A_7\rangle\,. \quad (35)$$

Setting $m_{\tilde g} = m_{\tilde b} \equiv \tilde m$ and $R_{\rm BR} = 1$, the gluino contributions to $\langle A_7\rangle$ read

$$\langle A_7\rangle = 2.5 \times 10^{-3} \left(\frac{500\,\mathrm{GeV}}{\tilde m}\right)^2 |\xi_L| \sin\gamma_L \left(1 + 2\frac{\mu\tan\beta - A_b}{\tilde m}\right)\,. \quad (36)$$

4.4. $S_{\phi K_S}$ and $S_{\eta' K_S}$

The expression for the mixing-induced CP asymmetries in B_d decays to final CP eigenstates f is

$$S_f = \sin(2\beta + \phi_\Delta + \delta_f)\,, \quad (37)$$

where ϕ_Δ is the new phase in $B_{d,s}$ mixing defined in (10). For the tree-level decay $f = \psi K_S$, $\delta_f = 0$, while for the penguin-induced modes $B \to \phi(\eta')K_S$, the contribution from the decay amplitude can be written as [17]

$$\delta_f = 2\arg\left(1 + a_f^u e^{i\delta} + \sum_{i\geq 3} b_{i,f}^c C_i^{\rm NP}\right) \quad (38)$$

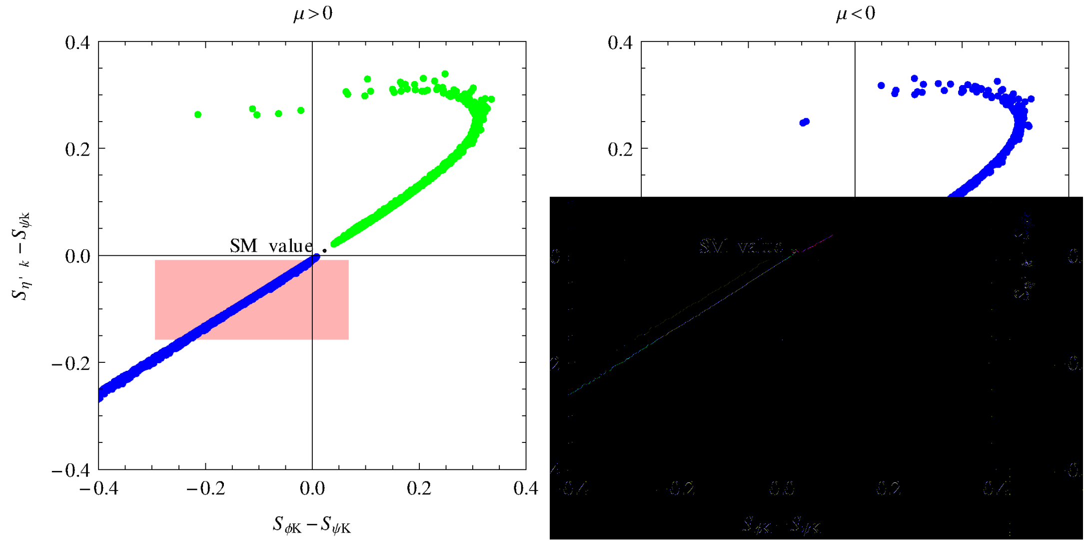

Figure 1: Correlation between $S_{\phi K_S} - S_{\psi K_S}$ and $S_{\eta' K_S} - S_{\psi K_S}$ for positive μ (left) and negative μ (right), showing points with $\gamma_L > 0$ (blue) and $\gamma_L < 0$ (green). The shaded region shows the 1σ experimental ranges.

where $\delta = \gamma_{\text{CKM}} = \phi_3$ is the usual CKM angle and the a_f^u and $b_{i,f}^c$ can be found in [17].

For the gluino contributions, setting $m_{\tilde{g}} = m_{\tilde{b}} \equiv \tilde{m}$, we obtain

$$\sum_{i=3}^{6} b_{i,\phi K_S}^c C_i^{\text{NP}} = -0.73 \times 10^{-2} \left(\frac{500\,\text{GeV}}{\tilde{m}}\right)^2 |\xi_L| e^{i\gamma_L}\,, \tag{39}$$

$$\sum_{i=3}^{6} b_{i,\eta' K_S}^c C_i^{\text{NP}} = -1.10 \times 10^{-2} \left(\frac{500\,\text{GeV}}{\tilde{m}}\right)^2 |\xi_L| e^{i\gamma_L}\,, \tag{40}$$

$$b_{8,\phi K_S}^c C_8^{\text{NP}} = -1.82 \times 10^{-2} \left(\frac{500\,\text{GeV}}{\tilde{m}}\right)^2 |\xi_L| e^{i\gamma_L} \left(1 + 2\frac{\mu \tan\beta - A_b}{\tilde{m}}\right), \tag{41}$$

$$b_{8,\eta' K_S}^c C_8^{\text{NP}} = -1.10 \times 10^{-2} \left(\frac{500\,\text{GeV}}{\tilde{m}}\right)^2 |\xi_L| e^{i\gamma_L} \left(1 + 2\frac{\mu \tan\beta - A_b}{\tilde{m}}\right). \tag{42}$$

The effects of the QCD and chromomagnetic penguins in the above expressions are comparable, with the exception of the left-right mixing piece only present for the chromomagnetic ones.

4.5. Numerical analysis

In figures 1 and 2 we show the correlations between the CP asymmetries, scanning the gluino mass between 0.5 and 1 TeV, the sbottom mass, the μ term and A_b between 0.2

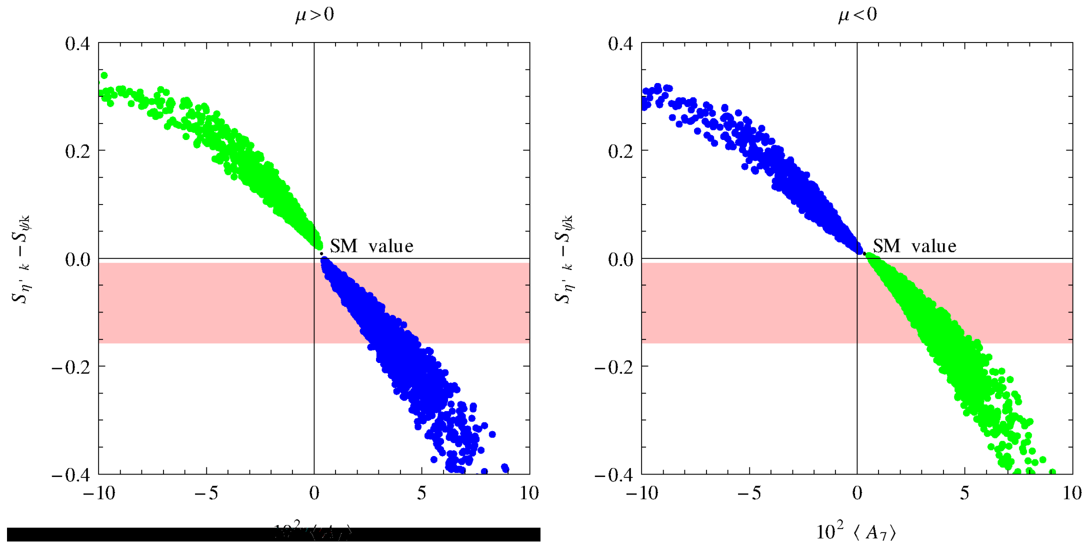

Figure 2: Correlation between $\langle A_7(B \to K^*\mu^+\mu^-)\rangle$ and the difference $S_{\eta' K_S} - S_{\psi K_S}$ for positive μ (left) and negative μ (right), showing points with $\gamma_L > 0$ (blue) and $\gamma_L < 0$ (green). The shaded region is the 1σ experimental range.

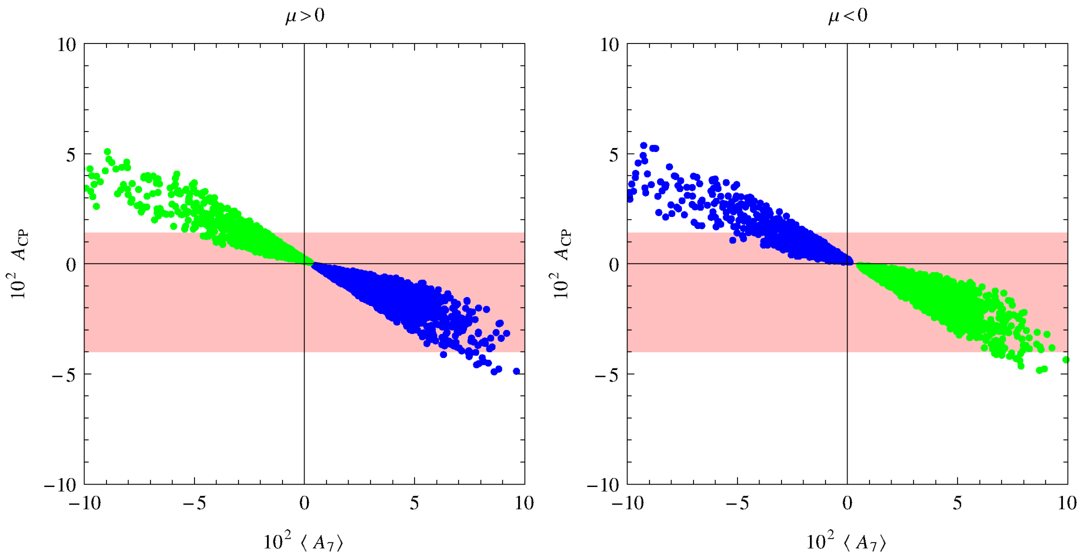

Figure 3: Correlation between $\langle A_7(B \to K^*\mu^+\mu^-)\rangle$ and the new physics contributions to the CP asymmetry in $B \to X_s\gamma$ for positive μ (left) and negative μ (right), showing points with $\gamma_L > 0$ (blue) and $\gamma_L < 0$ (green). The shaded region is the 1σ experimental range for $A_{\rm CP}(B \to X_s\gamma)$, valid for vanishing SM contribution (and otherwise subject to an appropriate shift).

and 0.5 TeV and $\tan\beta$ between 2 and 10. We require the $\Delta F = 2$ observables to be in the region where the CKM tensions are reduced (cf. (11–13)). The maximum size of the effects is mostly limited by the $\mathrm{BR}(B \to X_s\gamma)$ constraint, which we require to be fulfilled at the 3σ level including SM and gluino contributions only (cf. the discussion at the end of sec. 4.1).

Figure 1 shows the correlation between the mixing-induced CP asymmetries $S_{\phi K_S}$ and $S_{\eta' K_S}$ in relation to $S_{\psi K_S}$, effectively subtracting the contribution due to the modified B_d mixing phase. The experimental 1σ ranges corresponding to the average in table 1 are shown as shaded regions. Due to the $\tan\beta$ enhanced terms in (41, 42), large effects are easily possible. A negative value for these differences, as currently indicated by the central values of the measurements, can be obtained for $\gamma_L > 0$ ($\gamma_L < 0$) if $\mu > 0$ ($\mu < 0$). For a given sign of the μ term, the sign of the $\Delta B = 1$ CP asymmetries can thus serve to distinguish between the two solutions for the phase γ_L in (13) allowed by the $\Delta F = 2$ analysis.

Figure 2 shows the correlation between $S_{\eta' K_S}$ and the CP asymmetry $\langle A_7 \rangle$ in $B \to K^*\mu^+\mu^-$. Values up to $\pm 10\%$ would be attainable for $\langle A_7 \rangle$, while the current measurement of $S_{\eta' K_S}$ implies, at the 1σ level, $0 < \langle A_7 \rangle < 5\%$.

Finally, figure 3 shows the correlation between the CP asymmetry in $B \to X_s\gamma$ and the new physics contribution to $\langle A_7 \rangle$. $A_{\mathrm{CP}}(B \to X_s\gamma)_{\mathrm{NP}}$ attains values up to $\pm 5\%$. Imposing the 1σ experimental range allowed for $S_{\eta' K_S}$, this decreases to $-2\% < A_{\mathrm{CP}}(B \to X_s\gamma)_{\mathrm{NP}} < 0\%$. In the plots, we also show the 1σ experimental range for $A_{\mathrm{CP}}(B \to X_s\gamma)$ (cf. table 1), assuming a vanishing SM contribution. If the SM contribution is sizable due to long-distance effects (see section 4.2), the experimental constraints has to be shifted accordingly.

5. Discussion and conclusions

We have studied CP asymmetries in B decays in supersymmetry with a $U(2)^3$ flavour symmetry suggested in [10] as an alternative to Minimal Flavour Violation, addressing the SUSY flavour and CP problems, partially explaining the hierarchies in the Yukawa couplings and solving tensions in the CKM fit related to CP violation in meson mixing.

Even in the absence of flavour-blind phases, we find potentially sizable CP violating contributions to $\Delta B = 1$ decay amplitudes. We identify the dominant contributions to arise in the magnetic and chromomagnetic dipole operators due to their sensitivity to chirality violation, with subleading contributions in semi-leptonic and QCD penguin operators. The most promising observables are the mixing-induced CP asymmetries in non-leptonic penguin decays like $B \to \phi K_S$ or $B \to \eta' K_S$, angular CP asymmetries in $B \to K^*\mu^+\mu^-$, and the direct CP asymmetry in $B \to X_s\gamma$ (barring potential uncertainties in controlling the SM predictions in the radiative [15] and non-leptonic modes [17, 18]).

Due to the different dependence on the sparticle masses, we cannot predict a clear-cut correlation between CP violating $\Delta F = 1$ and $\Delta F = 2$ observables. However, we have demonstrated that observable effects in $\Delta F = 1$ CP asymmetries are certainly compatible with the pattern of deviations from the SM suggested by $\Delta F = 2$ observables, if interpreted in terms of this supersymmetric framework. Interestingly, while we considered a setup without flavour-blind phases, the correlations between $\Delta F = 1$ observables turn out to be very similar to those in MFV [25, 26] or in effective MFV [7] with flavour-blind phases. The main difference between the two cases are the CP violating effects in K and B mixing, which occur in $U(2)^3$, but not in (effective) MFV. We view this as an interesting example of the usefulness of correlated studies of $\Delta F = 1$ and $\Delta F = 2$ observables as a handle to distinguish between models. Such studies would become extremely interesting in presence of direct evidences of a hierarchical sparticle spectrum from the LHC.

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A. Loop functions

$$f_G(x) = \frac{2(73 - 134x + 37x^2)}{39(x-1)^3} - \frac{2(18 - 27x + x^3)}{13(x-1)^4}\ln x \qquad f_G(1) = 1$$

$$f_\gamma(x) = -\frac{2(2 - 7x + 11x^2)}{3(x-1)^3} + \frac{4x^3}{(x-1)^4}\ln x \qquad f_\gamma(1) = 1$$

$$f_7(x) = \frac{2(-1 + 5x + 2x^2)}{(x-1)^3} - \frac{12x^2}{(x-1)^4}\ln x \qquad f_7(1) = 1$$

$$g_7(x) = -\frac{6x(1 + 5x)}{(x-1)^3} + \frac{12x^2(2 + x)}{(x-1)^4}\ln x \qquad g_7(1) = 1$$

$$f_8(x) = \frac{-19 - 40x + 11x^2}{5(x-1)^3} - \frac{6x(-9 + x)}{5(x-1)^4}\ln x \qquad f_8(1) = 1$$

$$g_8(x) = \frac{12x(11 + x)}{5(x-1)^3} + \frac{6x(-9 - 16x + x^2)}{5(x-1)^4}\ln x \qquad g_8(1) = 1$$

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