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The MSSM with a softly broken U$(2)^3$ flavor symmetry

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In this article we review the phenomenological consequences of radiative flavor-violation (RFV) in the MSSM. In the model under consideration the $U(3)^3$ flavor symmetry of the gauge sector is broken in a first step to $U(2)^3$ by the top and bottom Yukawa couplings of the superpotential (and possibly also by the bilinear SUSY-breaking terms). In a second step the remaining $U(2)^3$ flavor symmetry is softly broken by the trilinear A-terms in order to obtain the measured quark masses and the CKM matrix of the Standard Model (SM) at low energies.

The phenomenological implications of this model depend on the actual choice of the SUSY breaking A-terms. If the CKM matrix is generated in the down sector (by A^d), $B_s \to \mu^+\mu^-$ receives non-decoupling contributions from Higgs penguins which become important already for moderate values of $\tan\beta$. Also the $B_s - \bar{B}_s$ mixing amplitude can be significantly modified compared to the SM prediction including a potential induction of a new CP-violating phase (which is not possible in the MSSM with MFV).

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1. Introduction

The most popular solution to the SUSY flavor problem is the hypothesis of minimal flavor violation (MFV) [1]: The soft SUSY-breaking terms are assumed to preserve a $U(3)_Q \times U(3)_u \times U(3)_d$ flavor symmetry, broken only by the Yukawa matrices $Y^{u\,(0)}$ and $Y^{d\,(0)}$. The imposed $U(3)^3$ symmetry is however in conflict with hints on new CP violating phases in $B_d - \bar{B}_d$ and $B_s - \bar{B}_s$ mixing, indicated by the D0 measurement of the dimuon charge asymmetry [2].[1] This effect cannot be explained within MSSM-MFV scenarios without enhancing $B_s \to \mu^+\mu^-$ far above the current experimental bounds[2]. Moreover, while a solution of the hierarchy problem favors stop masses well below 1 TeV, direct searches at the LHC resulted in tight bounds on the masses of the squarks of the first two generations. This conflict is a further challenge for the $U(3)^3$ symmetry.

These facts suggest to abandon the $U(3)^3$ flavor symmetry and to settle for a $U(2)^3$ for the first two generations in order to avoid conflicts with the tight constraints from Kaon and D-physics. A corresponding relaxed MFV scenario with the Yukawas $Y^{u\,(0)}$ and $Y^{d\,(0)}$ in the superpotential being the spurions breaking the $U(2)^3$ has been studied in Ref. [4]. We consider in this article a different scenario: We assume that the Yukawas $Y^{u\,(0)}$ and $Y^{d\,(0)}$ preserve the $U(2)^3$ flavor symmetry and that the soft SUSY-breaking trilinear A-terms A^u and A^d are the spurions breaking it. Such a model is quite appealing because it links the breaking of flavor symmetries to the breaking of supersymmetry. In the quark sector the $U(2)^3$ symmetry is then only softly broken, in the sense that the effective low-energy values of the Yukawa couplings $Y^u_{\rm eff}$ and $Y^d_{\rm eff}$, which are linked to the measured quark masses and CKM elements, are induced by A^u and A^d through radiative corrections. In this way the smallness of the light quark masses is explained via loop-suppression [5][3].

In this article we review our model of radiative flavor violation (RFV) and demonstrate that it can provide the above-mentioned new CP phase in $B_s - \bar{B}_s$ mixing in contrast to MFV. For a detailed study of further phenomenological consequences of the RFV scenario we refer to Ref. [6].

2. Radiative flavor violation (RFV)

In our scenario of RFV we assume that the bare Yukawa couplings $Y^{u\,(0)}$ and $Y^{d\,(0)}$ in the superpotential exhibit a $U(2)_Q \times U(2)_u \times U(2)_d$ flavor symmetry:

$$Y^{q(0)} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & y^q \end{pmatrix}, \qquad (q = u,d). \tag{2.1}$$

While the bilinear soft SUSY-breaking mass terms are assumed to possess the same symmetry, the trilinear A-terms A^u and A^d are the spurions breaking it. We perform $U(2)$-rotations on the left-

[1] Due to the recent LHCb measurement of the CP asymmetry in $B_s \to J/\psi\phi$ the situation is inconclusive at the moment. While the dimuon asymmetry measured by D0 points towards physics beyond the SM with $3.9\,\sigma$ significance and the CP asymmetry in $B_s \to J/\psi\phi$ measured by CDF has the same sign, LHCb obtained the opposite sign for the phase in $B_s - \bar{B}_s$ mixing (compatible with zero) from $B_s \to J/\psi\phi$.

[2] In MFV scenarios with an extended Higgs sectors, however, a large CP phase in $B - \bar{B}$ mixing is possible [3].

[3] For the corresponding analysis in the lepton sector see Ref. [7]

and right-handed up- and down-quark superfields to fix the basis in flavor space such that

$$A^q = \begin{pmatrix} A^q_{11} & 0 & A^q_{13} \\ 0 & A^q_{22} & A^q_{23} \\ A^q_{31} & A^q_{32} & A^q_{33} \end{pmatrix}, \qquad (q=u,d). \tag{2.2}$$

Note that the resulting basis is not a weak eigenbasis because left-handed up- and down-fields have to be rotated independently in order to diagonalize the 2×2 blocks of A^u and A^d simultaneously. Since in this basis no sources of flavor-violating (1,2) elements are present in the squark-mass matrices, the corresponding CKM matrix $V^{(0)}_{2\times 2}$ equals the Cabibbo matrix $V_{2\times 2}$ known from experiment (up to negligible corrections arising from loops involving a $1\to 3\to 2$ transition):

$$V^{(0)} = \begin{pmatrix} \cos\theta_C & \sin\theta_C & 0 \\ -\sin\theta_C & \cos\theta_C & 0 \\ 0 & 0 & 1 \end{pmatrix}. \tag{2.3}$$

The $U(2)^3$-breaking in the squark sector leaks into the quark sector via loop effects. The measured quark masses of the first two generations and the measured CKM elements are manifestations of the $U(2)^3$-breaking and as a consequence they must be directly related to the $A^{u,d}_{ij}$. Neglecting multiple flavor transitions (except for $1\to 2\to 3$ transitions) and small quark mass ratios one has

$$\begin{aligned} m_{q_i} &= a_q \frac{A^q_{ii}}{\mu_A} v_q\,, \qquad\qquad (q=u,d, \quad i=1,2)\,,\\ V_{cb} \approx -V^*_{ts} &= b_d \frac{A^d_{23}}{\mu_A}\frac{v_d}{m_b} - b_u \frac{A^u_{23}}{\mu_A}\frac{v_u}{m_t}\,,\\ V_{ub} &= b_d \left(\frac{A^d_{13}}{\mu_A} + V_{us}\frac{A^d_{23}}{\mu_A}\right)\frac{v_d}{m_b} - b_u \frac{A^u_{13}}{\mu_A}\frac{v_u}{m_t}\,,\\ -V^*_{td} &= b_d \frac{A^d_{13}}{\mu_A}\frac{v_d}{m_b} - b_u \left(\frac{A^u_{13}}{\mu_A} + V^*_{cd}\frac{A^u_{23}}{\mu_A}\right)\frac{v_u}{m_t}\,, \end{aligned} \tag{2.4}$$

where $\mu_A = \mathcal{O}(A^q_{ii})$ is a redundant mass scale introduced to render a_q, b_q dimensionless. The coefficients a_q, b_q are obtained by explicit evaluation of the self-energy diagrams inducing the quark mass terms. Restricting ourselves to SUSY-QCD contributions we find at first order in the mass insertion approximation

$$a_{u,d} = -\frac{2\alpha_s}{3\pi} m_{\tilde g}\,\mu_A\, C_0\left(m^2_{\tilde g}, m^2_{\tilde q_L}, m^2_{\tilde u_R,\tilde d_R}\right), \qquad b_{u,d} = -\frac{2\alpha_s}{3\pi} m_{\tilde g}\,\mu_A\, C_0\left(m^2_{\tilde g}, m^2_{\tilde q_L}, m^2_{\tilde t_R,\tilde b_R}\right). \tag{2.5}$$

Here $m_{\tilde q_L}$, $m_{\tilde u_R}$ and $m_{\tilde d_R}$ denote the common mass of the first two generations of left- and right-handed up- and down-type squarks, respectivly. If one assumes a $U(3)^3$ flavor symmetry for the bilinear squark mass terms (rather than only $U(2)^3$), one has $a_q = b_q$.

Eq. (2.4) implies that the A-terms A^q_{13} and A^q_{23} exhibit a similar hierarchy with respect to each other as the CKM-elements V_{ub} and V_{cb}, and in particular that $A^q_{11}/A^q_{22} = m_{q_1}/m_{q_2}$. The overall smallness, on the other hand, of the masses of the first two quark generations and of the off-diagonal CKM elements $V_{3i,i3}$ is explained by the loop suppression in a_q, b_q. The SUSY-flavor problem is restricted to the quantities $A^q_{31,32}$ which are not constrained from the measured CKM elements since their contributions are suppressed by small quark mass ratios.

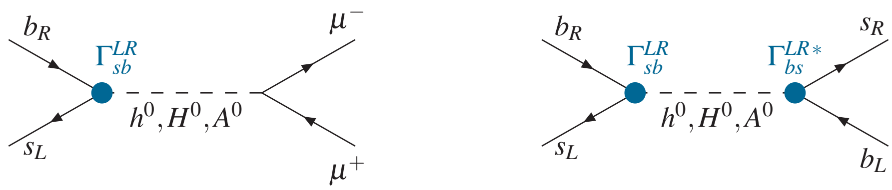

Figure 1: Higgs (double) penguin contributions to $B_s \to \mu^+\mu^-$ and to $B_s - \bar{B}_s$ mixing

3. Higgs (double) penguins: RFV vs. MFV

Loop diagrams involving the trilinear A-terms induce flavor-changing neutral Higgs-quark couplings $\Gamma_{q_i q_j}^{H_k^0\,LR} = \Gamma_{q_i q_j}^{LR} x_q^k$ $(q = u, d)$ where $H_k^0 = (H^0, h^0, A^0)$ [8]. The couplings Γ_{sb}^{LR}, Γ_{bs}^{LR} contribute via Higgs (double) penguin diagrams to the decay $B_s \to \mu^+\mu^-$ and to $B_s - \bar{B}_s$ mixing (see Fig. 1). For the relevant Wilson coefficients one has

$$\begin{aligned} B_s \to \mu^+\mu^- : &\qquad C_S \propto \Gamma_{sb}^{LR}, \qquad C_S' \propto \Gamma_{bs}^{LR*};\\ B_s - \bar{B}_s \text{ mixing} : &\qquad C_2^{LR} \propto \Gamma_{sb}^{LR}\,\Gamma_{bs}^{LR*}. \end{aligned} \tag{3.1}$$

Note that contributions to $B_s - \bar{B}_s$ mixing which are proportional to $(\Gamma_{sb}^{LR})^2$ or $(\Gamma_{bs}^{LR*})^2$ are strongly suppressed and thus negligible. This is a consequence of a Peccei-Quinn symmetry obeyed by the tree-level Yukawa-Lagrangian and the tree-level Higgs potential of the MSSM [9].

As the effective couplings Γ_{sb}^{LR} and Γ_{bs}^{LR} break the $U(2)^3$ flavor symmetry, they must be directly related to the corresponding spurions. In the MFV scenario we thus have[4]

$$\Gamma_{sb}^{LR} \propto \left(Y^u Y^{u\dagger} Y^d\right)_{23} \approx y_b y_t^2 V_{ts}^* V_{tb}\,, \qquad \Gamma_{sb}^{RL} = \Gamma_{bs}^{LR*} \propto \left(Y^u Y^{u\dagger} Y^d\right)_{32}^* \approx y_s y_t^2 V_{ts}^* V_{tb}\,. \tag{3.2}$$

Therefore Γ_{bs}^{LR} is suppressed compared to Γ_{sb}^{LR} by the small quark mass ratio m_s/m_b. Experimental bounds on $B_s \to \mu^+\mu^-$ constrain Γ_{sb}^{LR} and, because of the suppression of Γ_{bs}^{LR} with respect to Γ_{sb}^{LR}, they render Higgs double penguin effects in $B_s - \bar{B}_s$ mixing negligible [10].

In the RFV framework the spurions are given by the A terms. We consider here the limiting case in which A^u is flavor-diagonal so that the CKM elements in Eq. (2.4) are solely generated from the A_{i3}^d ("CKM generation in the down sector"). In this case one has[5]

$$\Gamma_{sb}^{LR} = \frac{A_{23}^d}{\mu_A} \propto V_{ts}^*\,, \qquad \Gamma_{sb}^{RL} = \Gamma_{bs}^{LR*} \propto \frac{A_{32}^{d*}}{\mu_A} \propto V_{32}^{R*}\,. \tag{3.3}$$

Here we have defined

$$V_{23}^R = -V_{32}^{R*} \equiv c\,\frac{A_{32}^{d*}}{\mu_A}\,\frac{v_d}{m_b}\,, \qquad c = -\frac{2\alpha_s}{3\pi} m_{\tilde{g}}\,\mu_A\,C_0\left(m_{\tilde{g}}^2, m_{\tilde{b}_L}^2, m_{\tilde{q}_R}^2\right) \tag{3.4}$$

with $m_{\tilde{b}_L,\tilde{t}_L}$ denoting the common mass of the left-handed sbottom and stop. The introduction of the quantity V_{32}^R simplifies the notation and allows for an easy comparison with the size of V_{ts}.

[4]Full expressions for the Higgs (double) penguin contributions to $B_s \to \mu^+\mu^-$ and $B_s - \bar{B}_s$ mixing in the MSSM with MFV including non-decoupling effects and complex phases for SUSY parameters can be found in Ref. [11].

[5]Full expressions for the effective Higgs couplings in RFV can be found in Ref. [6].

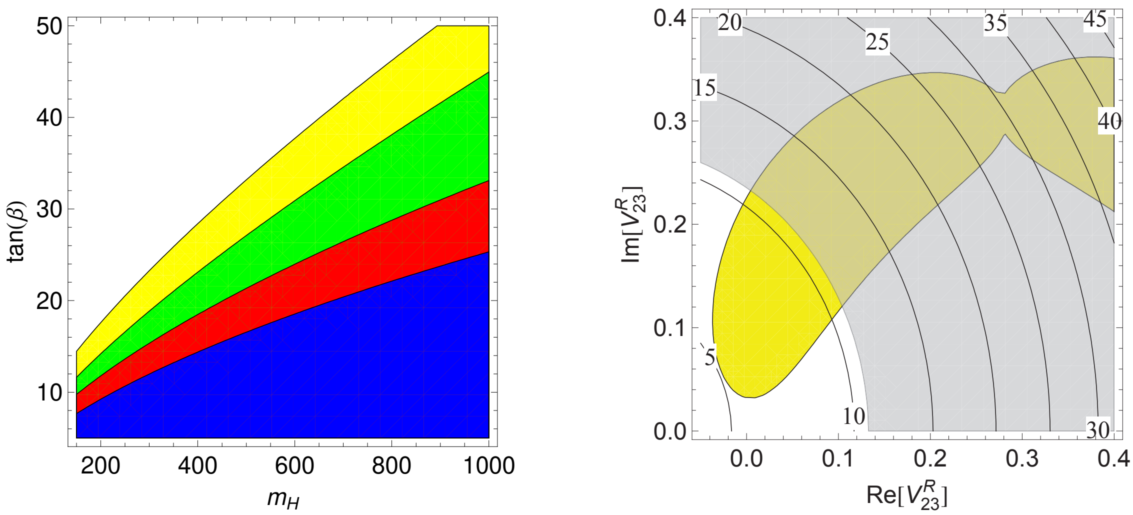

Figure 2: Left: Allowed region in the m_H–$\tan\beta$ plane for different values of ε_b from $\mathrm{Br}[\mathrm{B_s} \to \mu^+\mu^-] \leq 1.08 \cdot 10^{-8}$. Yellow: $\varepsilon_b = 0.005$, green: $\varepsilon_b = 0.01$, red: $\varepsilon_b = -0.005$, blue: $\varepsilon_b = -0.01$ (light to dark). Right: Correlation between $B_s \to \mu^+\mu^-$ and $B_s - \bar{B}_s$ mixing for $\varepsilon_b = 0.0075$, $m_H = 400$GeV for $\tan\beta = 11$. Yellow: Allowed region from $B_s - \bar{B}_s$ mixing (95% confidence level). The contour-lines show $\mathrm{Br}[\mathrm{B_s} \to \mu^+\mu^-] \times 10^9$. The grey area at the right side is excluded by the bound on $\mathrm{Br}[\mathrm{B_s} \to \mu^+\mu^-]$.

Since A^d_{32} is a free parameter of the theory, Γ^{LR}_{bs} is not related to Γ^{LR}_{sb} in RFV and in particular not suppressed with respect to the latter. Therefore Higgs double penguins can have sizable effects in $B_s - \bar{B}_s$ mixing, even in the light of present bounds on Γ^{LR}_{sb}, Γ^{LR}_{bs} from $B_s \to \mu^+\mu^-$.

In Fig. 2 on the left we show the allowed regions in the m_H–$\tan\beta$ plane from $B_s \to \mu^+\mu^-$ for different values of ε_b. On the right the correlation between $B_s - \bar{B}_s$ mixing and $B_s \to \mu^+\mu^-$ is shown for $m_H = 400$GeV, $\varepsilon_b = 0.0075$ and $\tan\beta = 11$. Note that there is a region in parameter space which can explain a potential new phase in $B_s - \bar{B}_s$ mixing and which is compatible with the current limits on $\mathrm{Br}[\mathrm{B_s} \to \mu^+\mu^-]$. Moreover, if the hints for a sizable new-physics contribution to $B_s - \bar{B}_s$ mixing persist, $B_s \to \mu^+\mu^-$ will necessarily be enhanced. LHCb will be able to probe this correlation in the near future.

If the CKM elements are generated from the A^u_{i3} terms (“CKM generation in the up sector”), interesting effects occur in the rare decays $K \to \pi\nu\bar{\nu}$ (see Ref. [6] for details).

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