FLAVOR VIOLATION AS A PROBE OF THE MSSM HIGGS SECTOR

CHRISTOPHER KOLDA

Theoretical Physics Group, Lawrence Berkeley National Laboratory
University of California, Berkeley, CA 94720, USA

At moderate to large $\tan \beta$, it is no longer possible to simultaneously diagonalize the masses of quarks and their couplings to the neutral Higgs bosons. The resulting flavor violations of the form $\bar{b}_R d_L \phi$ and $\bar{b}_R s_L \phi$ do not generate large meson-antimeson mixing amplitudes but do generate large contributes to rare decays such as $B_s \to \mu \mu$. Run II of the Tevatron will probe a large region of interesting MSSM parameter space through this decay channel. This talk is based on results obtained with K.S. Babu and presented in [1].

One of the most interesting facts about the minimal supersymmetric standard model (MSSM) is that it is not the supersymmetric minimal standard model. Holomorphy and anomaly cancellation force an additional Higgs doublet onto the model. Yet despite this extra Higgs doublet, it is still tempting to think of the minimal standard model as being the low-energy limit of the MSSM rather than its 2-Higgs doublet extension (the 2HDSM).

In a series of papers my coauthors and I have demonstrated that there are important new classes of phenomena associated with the 2HDSM-like limit of the MSSM which do not occur in a standard model-like limit. (Similar effects have been studied in [3].) These phenomena arise because the neutral Higgs bosons of the MSSM need not couple to fermions with strength proportional to the fermion mass, contrary to common folklore.

Consider first the non-supersymmetric 2HDSM. It is well known that such models must be carefully arranged in order to avoid large flavor-changing neutral currents (FCNCs). In particular, one must ensure that the couplings of the neutral Higgs states to fermions are proportional to the fermion masses. Were they not, Yukawa couplings would not be diagonal in the fermion mass eigenbases. The most economical method for guaranteeing this proportionality is to segregate the two Higgs doublets, allowing one to couple only to the $\bar{U}_R Q_L$ bilinear and the other to couple only to $\bar{D}_R Q_L$:

$$-\mathcal{L} = \bar{U}_R Y_{U} Q_L H_u + \bar{D}_R Y_{D} Q_L H_d + h.c.$$ (1)

where $U, D, Q$ are 3-vectors in flavor space and $Y_{U,D}$ are $3 \times 3$ matrices. Thus $H_u$ is alone responsible for giving mass to $u$-quarks and $H_d$ to $d$-quarks. When the Higgs fields mix to form a Higgs mass eigenstate $\phi = \{h^0, H^0, A^0\}$, the coupling to any given fermion will either come through the $H_u$ or the
$H_d$ component of $\phi$, but not both. Therefore $\phi$ also couples proportional to fermion mass, though with reduced strength.

This type of model, often called a “Type-II” model, is also stable against radiative corrections. Undesirable couplings of the form

$$- \mathcal{L}_{\text{bad}} = \bar{U} R Y U Q L H_u^* + \bar{D} R Y D Q L H_d^* + h.c.$$  \hfill (2)$$

(which would generate FCNCs) are forbidden by a $Z_2$ symmetry under which one of the two Higgs doublets is odd.

The MSSM is also a type-II model, but it possesses no such $Z_2$ symmetry. The superpotential of the MSSM is given by:

$$W = \hat{U} R Y U \hat{Q} L \hat{H}_u + \hat{D} R Y D \hat{Q} L \hat{H}_d + \mu \hat{H}_u \hat{H}_d$$  \hfill (3)$$

where I am ignoring leptons. It is the presence of a non-zero $\mu$-term which breaks the $Z_2$; and experimentally we know $\mu \neq 0$ due to the absence of a massless neutralinos in our colliders. However, the MSSM does not need any $Z_2$ to preserve the form of the Yukawa interactions. As long as SUSY is unbroken, holomorphy and the non-renormalization theorem guarantee that the dangerous Yukawa interactions of Eq. (2) cannot arise.

But SUSY is broken. And without any $Z_2$ to protect the form of (2), there is nothing to prevent couplings such as those in (2) from arising. In fact, they are known to arise after SUSY-breaking. In its present form, this observation was first made by Hall, Rattazzi and Sarid (HRS). They found that $d$-quarks can receive a 1-loop correction to their masses from the operator $\bar{D} R Q L H_u^*$. Though the coefficient of this new operator is small ($\sim \frac{1}{16\pi^2}$), the contribution of this operator to the $d$-quarks masses is enhanced by the ratio of the vev of $H_u$ to $H_d$, that is, by $\tan \beta$. Further, the leading diagrams do not decouple in the heavy SUSY mass limit, so long as all SUSY mass scales became large together.

The presence of a $\bar{D} R Q L H_u^*$ operator should be enough clue that something interesting will happen if we leave the full flavor structure of the HRS calculation intact. HRS considered two diagrams (shown in Fig. 1) which contribute to the $d$-quark masses, the larger mediated by gluinos, the smaller by charged Higgsinos. Both diagrams generate couplings between $\bar{D} R Q L$ and $H_u^*$. The gluino diagram, however, has a trivial flavor structure. Thus when we diagonalize the canonical $d$-quark mass term, we will simultaneously diagonalize this new contribution, preventing any flavor-violation. The second diagram is not so trivial thanks to the top squarks in the loop. Here the flavor structure of the new operator is given by $Y_D Y_U^T Y_U$, which is not diagonalized in the same basis as $Y_D$. Thus it is this second diagram which generates the flavor-changing. Simplify by diagonalizing the $u$-quark Yukawa interactions (similar effects are not sizable in the $u$-sector), leaving an effective...
Figure 1. Diagrams which contribute to $D_R Q_L H_u^*$ operator.

The Lagrangian in the $d$-sector:

$$\mathcal{L} = \bar{D}_R \mathbf{D} \mathbf{V}^0 \mathbf{Q}_L \mathbf{H}_d + \bar{D}_R \mathbf{D} \mathbf{V}^0 \mathbf{[} \epsilon_g + \epsilon_u \mathbf{U}^\dagger \mathbf{]} \mathbf{Q}_L \mathbf{H}_u^* + \text{h.c.} \quad (4)$$

where $\mathbf{U}, \mathbf{D}$ are the diagonalized Yukawa matrices; $\epsilon_g$ and $\epsilon_u$ are the contributions from the two diagrams in Fig. 1:

$$\epsilon_g \simeq (2\alpha_3 / 3\pi) \mu^* M_3, \quad \epsilon_u \simeq (1 / 16\pi^2) \mu^* A_U, \quad (5)$$

and $V^0$ is a unitary matrix which becomes the CKM matrix $V$ as $\epsilon_u \rightarrow 0$.

Now consider only the neutral current contributions. In the $u$-quark mass eigenbasis, the SU(2) partner of $U_L$ is not $D_L$ but rather $D'_L = V D_L$. Then the neutral current Lagrangian has the form:

$$\mathcal{L}_{NC} = \bar{D}_R \mathbf{D} \mathbf{V}^0 \mathbf{V} \mathbf{D}_L \mathbf{H}_d^0 + \bar{D}_R \mathbf{D} \mathbf{V}^0 \mathbf{[} \epsilon_g + \epsilon_u \mathbf{U}^\dagger \mathbf{]} \mathbf{V} \mathbf{D}_L \mathbf{H}_u^{0*} + \text{h.c.} \quad (6)$$

The matrix $V^0$ can be calculated explicitly since it must diagonalize the $d$-quark mass matrix: $V^0 \mathbf{Y}^\dagger \mathbf{Y} V^0 = \text{diag} (m_d^2, m_s^2, m_b^2)$ where

$$\mathbf{Y} = \mathbf{D} \mathbf{V}^0 \mathbf{[} 1 + \tan \beta (\epsilon_g + \epsilon_u \mathbf{U}^\dagger \mathbf{]}). \quad (7)$$

Keeping only the Yukawa couplings of the third generation, one can do the algebra, pulling out the flavor-changing pieces:

$$\mathcal{L}_{FCNC} = \frac{\bar{y}_b V^0_{tb}}{\sin \beta} \chi_{FC} \left[ V_{td} \bar{b}_R d_L + V_{ts} \bar{b}_R s_L \right] \left( \cos \beta H^0_{u} - \sin \beta H^0_{d} \right) + \text{h.c.} \quad (8)$$

where $\bar{y}_b = m_b / v_d$ ($\simeq 1$ for large $\tan \beta$) and $\chi_{FC} \simeq -\epsilon_u y_t^2 \tan \beta$. Note that $\chi_{FC}$ is proportional to $\epsilon_u$ as expected (complete expression in [1]). Of course, the $H_u$ and $H_d$ are not mass eigenstates, but it is simple to go to their mass eigenbasis so we will not show it here. Suffice to say that the flavor-changing coupling of the $h^0$ (lightest Higgs) goes to zero for large $m_A$, while those for the other Higgs bosons do not.

Where should we look for signs of this flavor changing? For most sources of FCNCs, it is meson–anti-meson mixing which provides the best opportunity for testing new effects. However here this is not the case. One can show that
there is a remnant flavor symmetry forbidding (at lowest order) $\Delta B = 2$
operators such as $B^0 - \bar{B}^0$ mixing. This left-over symmetry is present in
the MSSM only (not in generic two-Higgs models) and only at lowest order.
There are $\Delta B = 2$ contributions at higher order, but these have an additional
loop suppression and are therefore highly suppressed. Thus, it turns out that
meson–anti-meson mixing is a poor probe of this form of flavor-changing.

Rare $B$-decays ($\Delta B = 1$) provide another route, one which is not sup-
pressed by residual flavor symmetries. Since the mediation is via a Higgs,
final states are preferred according to their mass. Thus $B \to \tau\tau$ would be
ideal were it not so difficult experimentally. On the other hand, $B \to ee$
or $B \to X\nu\nu$ are highly suppressed. The optimal case is $B \to \mu\mu$ which was
studied in [1]. (Right behind is $B \to X\mu\mu$, only slightly weaker than the
purely leptonic case.) This decay can be probed at the Tevatron where $B_d$
and $B_s$ are produced in a ratio of about 3:1. However, the relative widths for
$B_{d,s} \to \mu\mu$ scale as $(V_{td}/V_{ts})^2 \sim 1/25$, so it is $B_s \to \mu\mu$ which will interest us.

For lack of space, only the final results will be shown here. Current CDF
limits of $\text{Br}(B_s \to \mu\mu) < 2 \times 10^{-6}$ constrain $m_{A^0} \lesssim 200$ GeV
given very large values of $\tan \beta \simeq 40 - 70$ and degenerate SUSY masses. At Run II,
that bound will increase to about 400 GeV after one year (1 fb$^{-1}$) and 650
GeV after several years (5 fb$^{-1}$). As long as $m_{A^0} \lesssim 1$ TeV and $\tan \beta$ is large,
the MSSM contributions to $B \to \mu\mu$ will exceed the SM prediction and be
observed before or at the LHC. And because these contributions persist even
for very heavy SUSY masses, the decay $B \to \mu\mu$ may be our first hint of
SUSY, long before SUSY partners are directly produced and observed.

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