5.3 Supersymmetry

5.3.1 CP Asymmetries in Supersymmetry

$>\!\!-$ C. Kolda $-\!\!<$

The Minimal Supersymmetric Standard Model (MSSM) provides an abundant opportunity for discovering New Physics in CP-violating and/or flavor-changing b quark processes. In fact, the most general version of the MSSM provides an over-abundance, with 12 masses, 30 mixing angles and 27 phases in the (s)quark sector, beyond those of the Standard Model. Of these, the LHC has only limited ability to go beyond measurements of the masses, leaving 57 parameters unconstrained, even after finding and studying SUSY at hadron colliders.

The other 57 parameters are not, however, wholly unconstrained. If one were to take all phases and angles $\mathcal{O}(1)$ and all masses $\mathcal{O}(\mathrm{TeV})$, the MSSM would make predictions for CP violation and FCNCs in the K sector far beyond those observed. One therefore expects some organizing principle to be at work in the MSSM, constraining the masses and/or phases and/or mixing angles in order to avoid phenomenological trouble. This is the so-called "SUSY flavor problem".

The source of the problem is that quarks get their masses only by electroweak symmetry breaking, while squarks get masses also by SUSY breaking. The SUSY-breaking contributions have no reason to be diagonal in the same bases as the SU(2)-breaking pieces, and so quark and squark mass eigenstates are not simultaneously defined. Unitarity of the mixing matrices is enough to force the quark-quark-gauge and squark-squark-gauge interactions to be diagonal (flavor-conserving), but the quark-squark-gaugino interaction will not be diagonal and will generate quark flavor-changing (and CP violation) through loops of squarks and gauginos.

There are three basic schemes which render the CP violation and FCNCs in the MSSM small: decoupling, alignment and degeneracy. Decoupling [147] is nothing more than the statement that if the MSSM sparticles are very heavy, then processes generated by them will be small. However, the masses required in order to actually get decoupling can be in the 100 to 1000 TeV range, far above the range where the MSSM plays an important role in electroweak symmetry breaking. Thus such models create their own mini-hierarchy problems. In the alignment scheme [148], one forces the squark and quark mass matrices to be diagonal in the same basis. However realistic alignment models are difficult to construct and often lead to large flavor-changing in the charm (D^0) sector.

The last of the suggestions is the one most often considered: degeneracy. If all squarks with the same gauge charges are degenerate in mass, then their contributions to flavor-changing/CP-violating processes exactly cancel. The degeneracy constraint is far more severe between the first and second generation squarks than with the third generation, because the constraints from the kaon system are so stringent. However, degeneracy is also more natural between the first and second generation, where Yukawa-induced renormalizations of the squark masses are small. While the current constraints for the third generation are much less severe, there is also reason to believe that some non-degeneracy is inevitable: the large top (and possibly bottom) Yukawa couplings will split the third generation off from the other two and will generate 1-3 and 2-3 squark mixings proportional to CKM elements. This can be seen simply by examination of the soft mass renormalization group equations. For example [149]:

$$\frac{d}{d\log Q}\left(m^2_{\tilde{Q}}\right)_{ij} \propto \left[Y_u^\dagger m^2_{\tilde{Q}} Y_u + Y_d m^2_{\tilde{Q}} Y_d^\dagger + \cdots\right]_{ij} . \tag{5.120}$$

In the basis in which Y_d is diagonal, the Y_u terms are rotated away from the diagonal by the CKM matrix:

$$(\delta m^2_{\tilde{Q}})_{ij} \simeq \frac{1}{8\pi^2}\log(M_X/M_{\mathrm{SUSY}}) \times (V^\dagger_{\mathrm{KM}} V_{\mathrm{KM}})_{ij} . \tag{5.121}$$

This is simply the most trivial example of physics that could cause the third generation to behave differently than the first two, and in fact this example introduces no new phases by itself (it is an example of minimal flavor violation [150]). But more complicated models exist, particularly those which attempt to explain the quark mass hierachy. So even though kaon physics may strongly constrain CP violation and FCNCs in the first two generations, there is still plenty of room for both in the third generation.

Why SUSY is Special

There is a hidden advantage to the scalar mass problems in SUSY. The lack of any strong flavor violation in kaons or B mesons seems to demand that the ultimate scale at which the usual Standard Model flavor problem (namely, why is $m_u \ll m_t$?) is resolved lie above 100 TeV or more. In fact there is no reason at all to prefer a scale for flavor physics near our current experimental sensitivity rather than far into the ultraviolet. But even if the flavor scale is far above the weak scale, SUSY may provide a unique window into this world, for two reasons. First, because SUSY is associated with the weak-scale hierarchy problem, its spectrum must lie near the weak scale. Thus the precision measurements at a Super B Factory will be sensitive to physics at the very scale at which SUSY is expected to be found.

But more importantly, it is the presence of so many scalar particles in SUSY that provides an extra sensitivity to high-scale flavor physics that would not normally be available. Scalar masses, through their renormalization, are sensitive to physics at all scales, from the weak scale to the far ultraviolet. In non-SUSY theories, quadratic divergences dominate this renormalization and it would not be clear how to interpret a scalar mass spectrum if one were observed. But in SUSY the scalar masses are only logarithmically renormalized, which allows the masses to be run up to high energies using the renormalization group. The presence of any non-trivial flavor physics anywhere below the SUSY-breaking scale tends to imprint itself on the spectrum of the scalar particles either through their renormalization group running or through threshold corrections at the flavor physics scale. In either case, flavor-violating operators which would normally be suppressed by powers of the flavor scale are instead suppressed only by powers of the SUSY mass scale (often with an additional large log enhancement); see Eq. (5.121). This idea has been particularly fruitful (at least theoretically) for probing the structure of the neutrino mass matrix and its correlations with $\tau \to \mu\gamma$ and $\mu \to e\gamma$. It is also the basic idea underlying several of the approaches [151] to $B \to \phi K_s^0$ that will be outlined in the next two sections.

Thus SUSY, which by itself provides no new insights into the question of flavor, may in fact be the mechanism by which we are finally able to gain experimental insights into flavor. It is for this reason that considerations of SUSY models and sensitivities will play an extremely important role in the future of high precision heavy flavor physics.

"Flavors" of SUSY

Because the MSSM requires some external organizing principle in order to keep the theory even remotely viable, the kinds of signals one expects at colliders depend sensitively on the organizing principle itself. In the simplest case in which degeneracy is enforced, all flavor violation is due to the CKM matrix. This is true even for the non-universal corrections generated by the renormalization group equations. Such models provide good examples of Minimal Flavor Violation [150] and one can refer to the section on MFV earlier in this chapter for a discussion of the relevant phenomenology.

However if the scale at which non-trivial flavor physics lies is below the scale at which SUSY is broken in the visible sector, evidence of the flavor physics should be imparted on the scalar spectrum in some way, even if suppressed. It would not be surprising to find that the strongest flavor violation among the scalars would occur where the Yukawa couplings are the greatest, namely in the third generation interactions. Thus a Super B Factory is the natural place to search for these effects.

It is customary (for ease of calculation) to work in a basis in which the quark masses are diagonal as are the quark-squark-gluino interactions. This forces the 6×6 squark mass matrix to remain non-diagonal. In the limit of approximate degeneracy (or approximate alignment), we interpret the diagonal elements of the mass-squared matrix to be the left- and right-handed squark masses, and the off-diagonal elements as mass insertions denoted $(\Delta^d_{ij})_{AB}$ where $i \neq j$ are generation indices ($i, j = 1 \ldots 3$) and A, B denote left(L) and right(R). We then define a mixing parameter:

$$(\delta^d_{ij})_{AB} = \frac{(\Delta^d_{ij})_{AB}}{\tilde{m}^2}\,, \tag{5.122}$$

where $\tilde{m}$ is a typical squark mass. Kaon physics constrains $(\delta^d_{12})_{AB}$ (for all AB) to be much, much smaller than one [152]. Experimental agreement of $B^0 - \overline{B}^0$ mixing with the Standard Model prediction likewise constrains

$(\delta^d_{13})_{AB}$ [153]. Compared to these cases, constraints on $(\delta^d_{23})_{AB}$ are relatively weak. Specifically, the LL and RR insertions can be $\mathcal{O}(1)$ while the LR and RL can be $\mathcal{O}(10^{-2})$ due to constraints from $b \to s\gamma$.

Appearance of a sizable $(\delta^d_{23})_{AB}$ will generate non- Standard Model $b \to s$ transitions, affecting branching ratios and asymmetries in a number of processes including $B \to \phi K^0_S$, $B \to X_s \ell\ell$, $B \to X_s\gamma$, $B \to \eta^{(\prime)} K^0_S$, $B \to K^+K^-K^0_S$, $B_s \to \ell\ell$, Δm_{B_s} and others. (We will assume that there is no large flavor violation in the 1-3 sector; such violation could enter $B^0 - \overline{B}^0$ mixing, and from there affect $B \to \psi K^0_S$.)

Of these, the CP-violating phase in $B \to \phi K^0_S$, namely $\beta_{\phi K}$, is of particular importance. $\beta_{\phi K}$ has been measured by *BABAR* and Belle to an accuracy around 10%. As of this writing, the *BABAR* and Belle experiments are in disagreement about whether or not there is an anomaly in the experimental data on $\beta_{\phi K}$ (the data is reported in terms of the oscillation parameter $S_{\phi K}$). Because of the hint that there might be an anomaly, many groups have conducted analyses of the $B \to \phi K^0_S$ in the context of SUSY [151, 154, 155]. Regardless, decays like $B \to \phi K^0_S$ and other $b \to s$ processes are key testing grounds for SUSY flavor physics.

$b \to s$ transitions in SUSY

The calculation of the short-distance SUSY contributions to $B \to \phi K^0_S$ is relatively straightforward. There are two classes of contributions which bear discussion, namely loops of charginos and loops of gluinos. Chargino loops contribute to the amplitude for $B \to \phi K^0_S$ with a structure that mimics the Standard Model. In particular, in models with minimal flavor violation, there is a SUSY contribution to the branching fraction for $B \to \phi K$ but not to the CP-violating asymmetries. If we extend minimal models to include arbitrary new phases (but not mixings) then the CP asymmetries can receive new contributions, but these are generally small. It may be possible to push $S_{\phi K}$ down to zero, but it is appears to be difficult to go any lower [156].

In models with arbitrary phases *and* mixings, the chargino contributions can be even larger, but now they are typically dwarfed by gluino contributions. The gluino contributions are absent in minimally flavor-violating models, but dominate in the case of general mixings and phases. Two types of gluino-mediated diagrams typically dominate the amplitudes for $B \to \phi K$: the chromomagnetic moment and gluonic penguins. (For details of these calculations, see Ref. [154]).

There are two questions of particular importance in examining the SUSY contributions to the CP asymmetries in $B \to \phi K^0_S$: can SUSY provide a large deviation from the Standard Model in $S_{\phi K}$, and what other observables would be correlated with a large deviation? In doing so, it is natural to consider four distinct cases or limits, with the understanding that a realistic model might contain elements of more than one case. Those cases are labelled by the chirality of the squark mixing: LL, RR, LR and RL, where the first letter labels the s squark and second the b squark.

Of the four cases, the LL insertion is particularly well motivated. In particular, one expects $(\delta^d_{23})_{LL} \sim V_{ts}$ even in models with minimal flavor violation. In models in which the SUSY breaking occurs at a high scale, the insertion can be enhanced by an additional large logarithm. The RR insertion is less motivated in minimal SUSY models, but is naturally generated in grand unified (GUT) models with large neutrino mixing [151]. In this case, the large mixing in the neutrino sector (which is contained in the $\overline{5}$ of SU(5)) is transmitted by GUT and renormalization effects to the right-handed down quark sector, which is also part of the same GUT representation.

The physics consequences of the LL and RR insertions are very similar to one another. In both cases, measurable deviations in $S_{\phi K}$ can be obtained. A sizable deviation in $S_{\phi K}$, however, requires large $(\delta^d_{23})_{LL,RR} \sim \mathcal{O}(1)$ and a relatively light SUSY spectrum. In order to obtain a negative $S_{\phi K}$ using an LL or RR insertion one requires gluinos with mass below 300 GeV, for example. The strongest external constraints on such large insertions and light masses come from direct searches for gluinos and from $b \to s\gamma$; the latter only constrains the $Re(\delta^d_{23})_{LL}$ to be greater than about -0.5, while providing no constraint on the RR insertion.

In order to determine that the New Physics in $S_{\phi K}$ would be coming from an LL or RR insertion, it must be correlated to other observables. Deviations in $S_{\phi K}$ are well correlated with deviations in $C_{\phi K}$: measurements of $S_{\phi K}$ below the

Table 5-11. *Correlated signatures for an observation of $S_{\phi K}$ much smaller than $S_{\psi K}$, assuming a single SUSY d-squark insertion of the type indicated. The $\pm$ signs represent the sign of the corresponding observable.*

	LL	RR	LR	RL
(δ^d_{23})	$O(1)$	$O(1)$	$O(10^{-2})$	$O(10^{-2})$
SUSY masses	$\lesssim 300$ GeV	$\lesssim 300$ GeV	$\lesssim$ TeV	$\lesssim$ TeV
$C_{\phi K}$	$-$, small	$-$, small	$-$, small	$-$, can be large
$\mathcal{B}(B \to \phi K)$	SM-like	SM-like	varies	varies
$A_{\rm CP}^{b\to s\gamma}$	$+$, few %	SM-like	$+$, $\mathcal{O}(10\%)$	SM-like
Δm_{B_s}	can be large	can be large	SM-like	SM-like

Standard Model expectation correlate to negative values of $C_{\phi K}$. (Note that the calculation of $C_{\phi K}$ is very sensitive to the techniques used for calculating the long distance effects; these correlations are found using the BBNS [157] method.) However the deviations in $C_{\phi K}$ are at most $\mathcal{O}(10\%)$ and so will require a much larger data sample such as that available at a Super B Factory. More striking is the correlation with Δm_{B_s}, the $B_s - -\overline{B}_s$ mass difference. Large deviations in $S_{\phi K}$ due to an LL or RR insertion correlate directly with very large mass differences, far outside the range that will be probed at Run II of the Tevatron. Mass differences of the order of 100 ps^{-1} are not atypical in models with large LL or RR insertions, making their experimental measurement very difficult.

Specific to an LL insertion (rather than an RR) will be deviations in the CP asymmetry in $b \to s\gamma$. Large negative deviations in $S_{\phi K}$ correlate cleanly with positive CP asymmetries of the order of a few percent. Measuring these asymmetries will require of order 10 ab^{-1} of data and so call for a Super B Factory.

The picture presented by the LR and RL insertions is quite different. First, the LR and RL insertions would generically be suppressed with respect to the LL and RR insertions, because they break SU(2) and must therefore scale as $M_W/M_{\rm SUSY}$. However they generate new contributions to the chromomagnetic operators which are enhanced by $M_{\rm SUSY}/m_q$ $(q = s, b)$ and are therefore very effective at generating large deviations in $S_{\phi K}$. The LR insertion is the more well motivated, since one expects $\tilde{s}_L$–$\tilde{b}_R$ mixing to be proportional to the bottom Yukawa coupling, while $\tilde{s}_R$–$\tilde{b}_L$ mixing would come from the much smaller strange Yukawa. However it is possible to build reasonable flavor models in which this assumed hierarchy is not preserved and sizable RL insertions are generated [154].

In either case, whether LR or RL, strong constraints from the branching ratio of $b \to s\gamma$ force $(\delta^d_{23})_{LR,RL}$ to be $\mathcal{O}(10^{-2})$. Neither insertion generates an observable shift in Δm_{B_s}, but both can generate large shifts in the branching fraction for $B \to \phi K$. Of more interest are the correlations between $S_{\phi K}$, $C_{\phi K}$ and the CP asymmetry in $b \to s\gamma$. For measurements of $S_{\phi K}$ below the Standard Model , the LR insertion always predicts a negative $C_{\phi K}$, with values down to -0.3 when $S_{\phi K}$ goes as low as -0.6. On the other hand, negative contributions to $S_{\phi K}$ are associated with positive asymmetries in $b \to s\gamma$, often as large as 5% to 15%. These large asymmetries are a clear sign of the presence of an LR insertion, as opposed to LL insertions which give asymmetries of only a few percent.

For the RL case, the phenomenology is much the same except: (1) the values of $C_{\phi K}$ implied by a down deviation in $S_{\phi K}$ are even more negative, all the way down to $C_{\phi K} = -1$; (2) though the RL operators contribute to $b \to s\gamma$, they do not interfere with the Standard Model contribution and thus do not generate any new source of CP violation. Thus RL insertions predict no new observable CP asymmetry in $b \to s\gamma$.

Table 5-11 summarizes the correlations for each type of insertion. Of course, more than one insertion may be present, so one could generate large deviations in $S_{\phi K}$ with an LR insertion and large Δm_{B_s} with an LL insertion. Such combinations can be read off of the table.

In conclusion, observation of a significant (or any) deviation in the CP asymmetries in $B \to \phi K_S^0$ could be an early and strong indication of SUSY flavor physics. But it is the correlations between the ϕK_S^0 signal and other observables that will lead us to a deeper understanding of flavor in the MSSM.

5.3.2 SUSY at the Super B Factory

$\succ$– T. Goto, Y. Okada, Y. Shimizu, T. Shindou, and M. Tanaka –$\prec$

The Unitarity triangle and rare decays in three SUSY models

Among various candidates of physics beyond the Standard Model, SUSY is regraded as the most attractive possibility. The weak scale SUSY provides a solution of the hierarchy problem in the Standard Model. Although an extreme fine-tuning is necessary to keep the weak scale very small compared to the Planck scale within the Standard Model , SUSY theory does not have this problem, because of the cancellation of the quadratic divergence in the scalar mass renormalization. SUSY has attracted much attention since early 1990's, when three gauge coupling constants determined at LEP and SLC turned out to be consistent with the coupling unification predicted in SUSY GUT.

One of main motivations of the LHC experiment is a direct search for SUSY particles. The mass reach of colored SUSY particles will be about 2 TeV, an order-of-magnitude improvement from the present limit. It is quite likely that some signal of SUSY can be obtained in the early stage of the LHC experiment. It is therefore important to clarify the role of SUSY studies at a Super B Factory in the LHC era.

In order to illustrate how B physics can provide useful information to distinguish various SUSY models, we study SUSY effects in the length and angle measurements of the unitarity triangle and rare B decays for the following four cases in three SUSY models [158, 159, 99].

- Minimal supergravity model
- SU(5) SUSY GUT with right-handed neutrinos: Case 1 (degenerate case)
- SU(5) SUSY GUT with right-handed neutrinos: Case 2 (non-degenerate case)
- MSSM with a U(2) flavor symmetry

In the first model, all squarks and sleptons are assumed to be degenerate at a high energy scale such as the Planck scale, where SUSY breaking effects are transmitted to the observable sector from the hidden sector by the gravitational interaction. Flavor mixings and mass-splittings are induced by renormalization effects due to the ordinary quark Yukawa coupling constants, especially from a large top Yukawa coupling constant. The matrices which diagonalize the resulting squark mass matrices are approximately given by the CKM matrix, because this is the only source of the flavor mixing in the quark and squark sectors. In this sense, this model is a realization of so called "minimal flavor violation" scenario. We can consider new SUSY CP phases for the trilinear scalar coupling (A parameter) and the higgsino mass term (μ parameter). There phases are, however constrained by various electric dipole moment (EDM) experiments in the context of the minimal supergravity model, so that effects on B physics are relatively small [160].

A SUSY GUT with right-handed neutrinos is a well-motivated candidate of the physics beyond the Standard Model. Here, we consider an SU(5) SUSY GUT model and incorporate the seesaw mechanism for neutrino mass generation by introducing an SU(5) singlet right-handed neutrinos. In this model, large flavor mixing in the neutrino sector can affect the flavor mixing in the squark sector through renormalization of sfermion mass matrices [161, 162, 163, 164]. Since the lepton weak doublet and the right-handed down-type quark are included in the same SU(5) multiplet, the renormalization induces the flavor mixing in the right-handed down-type squark sector. At the same time, lepton flavor violation (LFV) in the charged lepton sector is also induced.

We consider two specific cases for the right-handed neutrino mass matrix, since constraints from LFV processes, especially from the $\mu \to e\gamma$ process, depend on the matrix significantly. From the seesaw relation, the light neutrino mass matrix is given by $m_\nu = y_\nu^T (M_N)^{-1} y_\nu (v^2 \sin^2 \beta/2)$ in the basis where the charged Higgs Yukawa coupling is diagonal. Here, y_ν, M_N, and β are the neutrino Yukawa coupling constant, the right-handed neutrino mass matrix, and the vacuum angle, respectively. On the other hand, the LFV mass term for the left-handed slepton are given by $(\delta m_L^2)^{ij} \simeq -(y_\nu^\dagger y_\nu)^{ij}(3m_0^2 + |A_0|^2) \ln{(M_P/M_R)}/8\pi^2$, where m_0, A_0, M_P, and M_R are the universal scalar mass, the trilinear scalar coupling, the Planck mass, and the right-handed mass scale. The first case is a degenerate case, where the right-handed neutrino mass matrix is proportional to a unit matrix. The off-diagonal element of the slepton mass matrix is related to the neutrino mixing matrix in this case, and therefore the large mixing angle suggested by the solar neutrino observation indicates a severe constraint on SUSY parameters from the $\mu \to e\gamma$ branching ratio. On the other hand, the constraint is relaxed, if we arrange the right-handed neutrino mass matrix such that the 1-2 and 1-3 mixings of $y^\dagger y$ vanish. We call this limiting case a non-degenerate case. Since the constraint to the SUSY parameter space is quite different in two cases, we calculate various observable quantities for both, and compare their phenomenological implications. Note that in the viewpoint of the LFV constraint the degenerate case represents a more generic situation than the non-degenerate case, because the $\mu \to e\gamma$ process generally puts a severe restriction to allowed ranges of SUSY parameters.

The minimal supersymmetric Standard Model (MSSM) with U(2) flavor symmetry was proposed sometime ago as a solution of the flavor problem in a general SUSY model [165, 166]. Unless the squark and slepton masses are in the multi-TeV range, there should be some suppression mechanism for flavor changing neutral current (FCNC) processes, especially for the squarks and sleptons of the first two generations. If we introduce a U(2) flavor symmetry, under which the first two generations are assigned to be doublets and the third generation is a singlet, we can explain the realistic pattern of the quark mass and suppress the unwanted FCNC at least for the kaon sector. FCNC of the bottom sector, on the other hand, can be interesting signals. We follow a specific model of this type according to Ref. [166]. In this model, there are many $\mathcal{O}(1)$ parameters in squark mass matrices, which we have scanned in a reasonable range.

We have calculated the following quantities for the above four cases in the three models.

- CP violation parameter ϵ_K in the $K^0 - \overline{K^0}$ mixing.
- $B_d - \overline{B}_d$ mixing and $B_s - \overline{B}_s$ mixing.
- The mixing induced CP violations in $B \to J/\psi\, K_S^0$ and $B \to \phi K_S^0$ modes.
- The mixing-induced CP violation in $B \to M_s\gamma$, where M_s is a CP eigenstate with a strange quark such as $K^*(\to K_S^0\pi^0)$.
- Direct CP violation in the inclusive $b \to s\gamma$ process.

These quantities provide several independent methods to look for New Physics. New Physics contributions in the mixing quantities may be identified from the consistency test of the unitarity triangle. The difference of the CP asymmetries in $B \to J/\psi\, K_S^0$ and $B \to \phi K_S^0$ implies existence of a new CP phase in the $b \to s$ transition amplitude. For the $b \to s\gamma$ process, a sizable direct CP asymmetry means a new phase in the $b \to s\gamma$ amplitude, while the mixing-induced asymmetry arises from the interference between the amplitudes with $b \to s\gamma_L$ and $\overline{b} \to \overline{s}\gamma_L$. Although this is suppressed by m_s/m_b in the Standard Model, New Physics effects can generate $O(1)$ asymmetry, if there is a $b \to s\gamma$ amplitude with the opposite chirality. Detailed description of our calculation is given in Ref. [158, 99].

The correlation among the CP asymmetry of the $B \to J/\psi\, K_S^0$ mode, the phase of V_{ub}^* element (ϕ_3), and the ratio of the $B_s - \overline{B}_s$ mixing and $B_d - \overline{B}_d$ = mixing ($\Delta m(B_s)/\Delta m(B_d)$) is shown in Fig. 5-30. In this figure, we have taken into account theoretical uncertainties due to the kaon bag parameters ($\pm 15\%$) and $f_B\sqrt{B_d}$ ($\pm 20\%$) and take $|V_{ub}/V_{cb}| = 0.09 \pm 0.01$. In the calculation, we have imposed various phenomenological constraints to restrict SUSY parameter space. These includes constraints from the Higgs boson and SUSY particle searches in collider experiments, the branching ratio of the $b \to s\gamma$ process, and various EDM experiments. We updated the previous calculation given

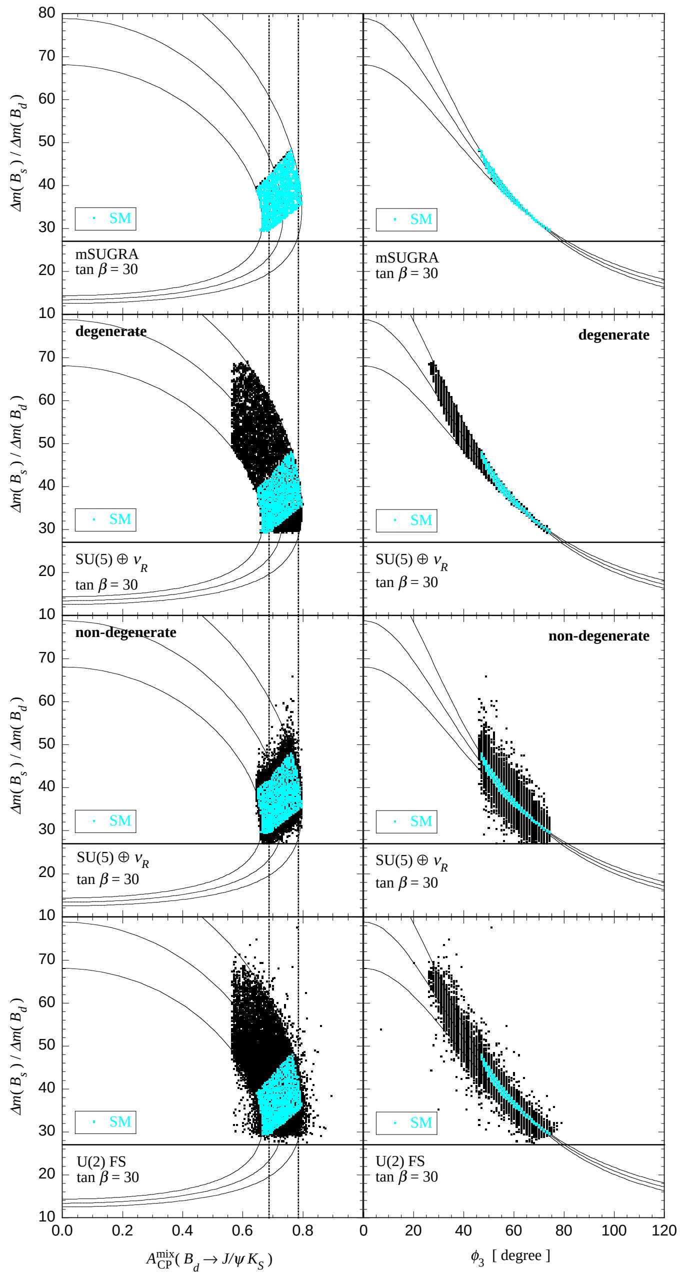

Figure 5-30. *$\Delta m(B_s)/\Delta m(B_d)$ versus the mixing-induced CP asymmetry of $B_d \to J/\psi K_S^0$ and ϕ_3 in the minimal supergravity model, SU(5) SUSY GUT with right-handed neutrinos for the degenerate and non-degenerate cases of the right-handed neutrino mass matrix, and the MSSM with a U(2) flavor symmetry. The light-colored regions show the allowed region in the Standard Model . The curves show the Standard Model values with $|V_{ub}/V_{cb}| = 0.08$, 0.09 and 0.10.*

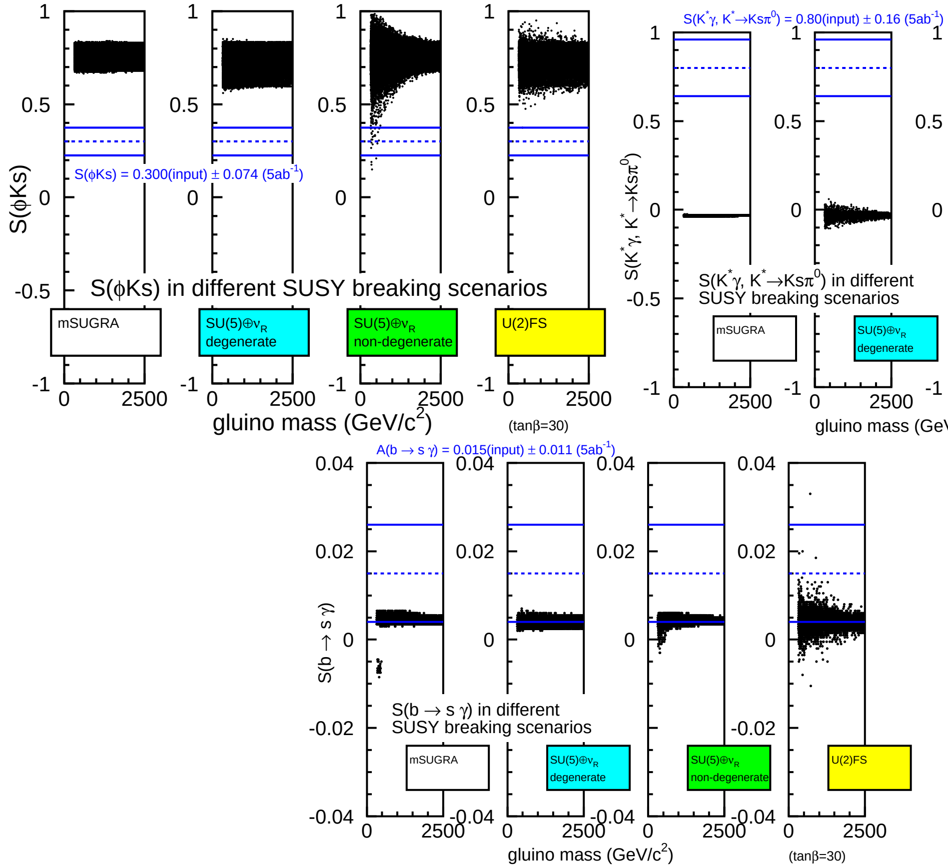

Figure 5-31. *Mixing-induced CP asymmetry in ϕK_S^0 and $M_s\gamma$ modes and direct CP asymmetry in $b \to s\gamma$ as a function of the gluino mass.*

in Ref. [158, 99] by taking account of the constraint on parameter space from the strange quark EDM contribution to the Hg EDM, which was pointed out recently [167]. For the two cases of the SU(5) GUT with right-handed neutrinos, we included the $\mu \to e\gamma$ constraint, which is especially important in the degenerate case. The figure corresponds to $\tan\beta = 30$. For neutrino parameters in the GUT model, we take a hierarchal light neutrino mass spectrum with the large mixing angle MSW solution. The right-handed neutrino masses are taken to be 4×10^{13} GeV for the degenerate case, and $5.7, 18, 45 \times 10^{13}$ GeV for the non-degenerate case.

We can see that possible deviations from the Standard Model prediction are small for the minimal supergravity model. In the SU(5) GUT with right-handed neutrinos, the pattern of the deviation is different for the two cases. For the degenerate case, $\Delta m(B_s)/\Delta m(B_d)$ can be enhanced from the Standard Model prediction, while the correct value of ϕ_3 is smaller than that expected in the Standard Model; this deviation, in fact, arises from a large SUSY contribution to ϵ_K. The deviation may become clear when the value of ϕ_3 is precisely determined from CP asymmetries of tree precesses such as $B \to DK$. For the non-degenerate case, the deviation can be seen only for $\Delta m(B_s)/\Delta m(B_d)$. We can conclude that SUSY contributions are large for the 1-2 generation mixing in the former case, and the 2-3 generation mixing in the latter case. More general type of deviations is possible for the MSSM with a U(2) flavor symmetry, because all three mixing diagrams can have large contributions.

The mixing-induced CP asymmetries in the $B \to \phi K_S^0$ and $B \to M_s\gamma$ modes and the direct CP asymmetry of $b \to s\gamma$ are shown in Fig. 5-31 for the four cases in the three models. Possible values of these observable quantities are plotted in terms of the gluino mass for $\tan\beta = 30$. These figures are updated from those in Ref. [99], taking into account the

Table 5-12. *Pattern of the deviation from the Standard Model predictions for unitarity triangle and rare decays. "$\surd$" means that the deviation can be large and "-" means a small deviation. "closed" in the first row of the B_d unitarity means that the unitarity triangle is closed among observables related to B_d, and the second and the third rows show whether deviation is observed from consistency check between the B_d unitarity and ϵ_K and $\Delta m(B_s)/\Delta m(B_d)$, respectively.*

	B_d unitarity			Rare Decays		
	closure	$+\epsilon_K$	$+\Delta m(B_s)$	$A^{\rm mix}_{CP}(B \to \phi K^0_S)$	$A^{\rm mix}_{CP}(B \to M_s\gamma)$	$A^{\rm dir}_{CP}(B \to X_s\gamma)$
mSUGRA	closed	-	-	-	-	-
SU(5) SUSY GUT (degenerate RHN)	closed	$\surd$	-	-	-	-
SU(5) SUSY GUT (non-deg. RHN)	closed	-	$\surd$	$\surd$	$\surd$	-
MSSM with U(2)	$\surd$	$\surd$	$\surd$	$\surd$	$\surd$	$\surd$

Hg EDM constant. The expected experimental sensitivities at a Super B Factory with integrated luminosity of 5 ab^{-1} are also indicated based on the study for the Super KEKB LoI. The central values are chosen for illustrative purpose. We can see that SUSY effects are large for the mixing-induced CP asymmetries for $B \to \phi K^0_S$ and $B \to M_s\gamma$ in the non-degenerate case of SU(5) SUSY GUT with right-handed neutrinos, whereas the corresponding deviations are small for the degenerate case. In the degenerate case the constraint from the $\mu \to e\gamma$ branching ratio is so strong that the effect in the 2-3 generation mixing is not sizable. In the non-degenerate case, the contribution from the sbottom-sdown mixing induces large effects in the $b \to s$ transition, because the $\mu \to e\gamma$ constraint is somewhat relaxed. For the U(2) case, we can see that all three deviations can be sizable.

Possible deviations from the Standard Model prediction in the consistency test of the unitarity triangle and rare decays are summarized in Table 5-12. The patterns of the deviations are different for these cases. For instance, observables related to the B_d unitarity triangle, namely $\Delta m(B_d)$, $|V_{ub}|$, ϕ_1 from the $B \to J/\psi_S$ mode, and ϕ_3 from the $B \to DK$ mode are consistent with a single triangle for the first tree cases in the table, but deviation can be observed if we compare ϵ_K and $\Delta m(B_s)/\Delta m(B_d)$ with the Standard Model prediction for the second and third cases. The deviation patterns are also different for various rare decay observables. These features are useful to distinguish different SUSY models at a Super B Factory.

B physics signals in the Snowmass Points & Slopes

It is expected that LHC experiments can significantly improve the search limit of SUSY particles. In a typical scenario like the minimal supergravity model, squarks and gluino will be found if their masses are below 2 TeV. Snowmass Points and Slopes (SPS) are proposed sets of benchmark parameters of SUSY parameter space [168]. Such model points and lines are selected as representative cases for phenomenological studies of SUSY theory, especially for SUSY particle searches in future collider experiments.

From the viewpoint of a Super B Factory, it is interesting to study possible flavor physics signals in these benchmark scenarios, and compare them with collider signals. In order to illustrate how LHC and a Super B Factory can be complementary to each other, we calculated FCNC processes and rare decays along several benchmark parameter lines for the two cases of SU(5) GUT with right-handed neutrinos. We should note that the benchmark points are mainly intended to select representative SUSY mass spectrum for physics analysis at collider experiments, whereas the flavor physics depends on how flavor off-diagonal terms in the squark/slepton mass matrices are generated. It is therefore conceivable that B physics can distinguish different SUSY models, even if the SUSY spectrum looks very similar.

We consider the following model-parameter lines, corresponding to four cases in the SPS list. These lines are defined by input parameters of the minimal supergravity model, namely the universal scalar mass (m_0), the gaugino mass

($m_{1/2}$), the universal trilinear coupling (A_0), and the vacuum ratio ($\tan\beta$). The sign of the higgsino mass term (μ) is taken to be positive.

- SPS 1a: $m_0 = -A_0 = 0.4 m_{1/2}, \tan\beta = 10$
- SPS 1b: $m_0 = 0.5 m_{1/2}, A_0 = 0, \tan\beta = 30$
- SPS 2: $m_0 = 2m_{1/2} + 850\mathrm{GeV}, \mathrm{A}_0 = 0, \tan\beta = 10$
- SPS 3: $m_0 = 0.25 m_{1/2} - 10\mathrm{GeV}, \mathrm{A}_0 = 0, \tan\beta = 10$

The lines are defined by varying $m_{1/2}$ The first two cases represent typical parameter points in the minimal supergravity model. (SPS 2b was only defined for a point with $m_{1/2} = 400$ GeV in [168], but here we generalize it to a line by varying $m_{1/2}$.) SPS 2 corresponds to the focus point scenario, where squarks and sleptons are rather heavy [169]. SPS 3 is a line in the co-annihilation region, where a rapid co-annihilation between a lighter stau and a LSP neutralino allows acceptable relic abundance for LSP dark matter. We take these input SUSY parameters for the SUSY GUT model, although the precise mass spectrum is not exactly the same as the minimal supergravity case due to additional renormalization effects from the neutrino Yukawa coupling, *etc.*

The results of the calculation for $\epsilon_K/(\epsilon_K)_{SM}$, $\Delta m(B_s)/(\Delta m(B_s))_{SM}$, $A^{\rm mix}_{CP}(B\to\phi K^0_S)$, and $A^{\rm mix}_{CP}(B\to M_s\gamma)$ are summarized in Table 5-13. In this calculation, we take the right-handed neutrino mass scale around 10^{14} GeV as before, and new phases associated with GUT interactions are varied. The calculation procedure is the same as that in Ref. [158, 99]. The table lists magnitudes of maximal deviations from the Standard Model prediction for each quantity. We do not list $A^{\rm dir}_{CP}(B\to X_s\gamma)$, because possible deviations are not large even in more general parameter space as described before. We see that the only sizable deviation appears for $\epsilon_K/(\epsilon_K)_{StandardModel}$ in the degenerate case of SPS 2. For other cases, it is difficult to distinguish these models from the prediction of the Standard Model or the minimal supergravity model with an integrated

Table 5-13. *Possible deviation from the Standard Model prediction for various observable quantities for benchmark parameter lines in the SU(5) SUSY GUT with right-handed neutrinos. The degenerate and non-degenerate cases for right-handed neutrinos are shown separately. SPS 1a, 1b, 2, and 3 are model-parameter lines defined in the text. The right-handed neutrino mass scale is taken to be $O(10^{14})$ GeV, and GUT phases are varied.*

Degenerate case	$\epsilon_K/(\epsilon_K)_{SM}$	$\Delta m(B_s)/(\Delta m(B_s))_{SM}$	$A^{\rm mix}_{CP}(B\to\phi K^0_S)$	$A^{\rm mix}_{CP}(B\to M_s\gamma)$
SPS 1a	$\lesssim$ 10%	$\lesssim$ 2%	$\lesssim$ 0.1%	$\lesssim$ 0.1%
SPS 1b	$\lesssim$ 10%	$\lesssim$ 2%	$\lesssim$ 0.5%	$\lesssim$ 0.2%
SPS 2	$\lesssim$ 100%	$\lesssim$ 2%	$\lesssim$ 0.3%	$\lesssim$ 0.5%
SPS 3	$\lesssim$ 5%	$\lesssim$ 2%	$\lesssim$ 0.1%	$\lesssim$ 0.1%
Non-degenerate case				
SPS 1a	$\lesssim$ 2%	$\lesssim$ 2%	$\lesssim$ 0.1%	$\lesssim$ 1%
SPS 1b	$\lesssim$ 1%	$\lesssim$ 2%	$\lesssim$ 0.5%	$\lesssim$ 2%
SPS 2	$\lesssim$ 1%	$\lesssim$ 3%	$\lesssim$ 1%	$\lesssim$ 3%
SPS 3	$\lesssim$ 2%	$\lesssim$ 2%	$\lesssim$ 0.1%	$\lesssim$ 1%

luminosity of 5 ab^{-1}. $\epsilon_K/(\epsilon_K)_{SM}$ is shown for the degenerate case as a function of the gluino mass in Fig. 5-32. For the case of SPS 2, the deviation of this size can be distinguished at a Super B Factory by improved measurements of quantities related to the unitarity triangle, especially ϕ_3. On the other hand, the $b\to s$ transition processes do not show large deviations even for the non-degenerate case for the selected model-lines. This is in contrast to the scatter plot in more general parameter space. We find that a large deviation occurs only for large values of the A_0 parameter, but the benchmark lines do not correspond to such cases. We should also note that a sizable deviation in the SPS 2 case can be seen even for a relatively heavy SUSY spectrum where squarks are 1 -2 TeV, which can be close to the discovery limit of SUSY at the LHC experiments.

In summary, we studied SUSY effects to various FCNC processes related to the unitarity triangle and rare B decay processes with a $b \to s$ transition. We considered the minimal supergravity model, two cases for the SU(5) SUSY GUT with right-handed neutrinos, and the MSSM with a U(2) flavor symmetry. We found that large deviations are possible in observable quantities with either 1-2 or 2-3 generation transition depending on the choice of the right-handed neutrino mass matrices and the neutrino Yukawa coupling constants in the GUT model. Various New Physics signals are possible in the U(2) model, while the deviation is small for the minimal supergravity model. These features are useful to identify possible SUSY models at a Super B Factory . We also consider SUSY parameter space based on benchmark scenarios of SPS. We observe that SUSY contribution can be large in ϵ_K for the case of SPS 2 (focus point scenario) with the degenerate right-handed neutrinos in SU(5) SUSY GUT. This example illustrates that a Super B Factory can provide important insight to the flavor structure of SUSY theory, which is complementary to what will be obtained at energy frontier collider experiments.

5.3.3 Electric Dipole Moment for ^{199}Hg atom and $B \to \phi K^0_s$ in Supersymmetric Models with Right-Handed Squark Mixing

>– J. Hisano and Y. Shimizu –<

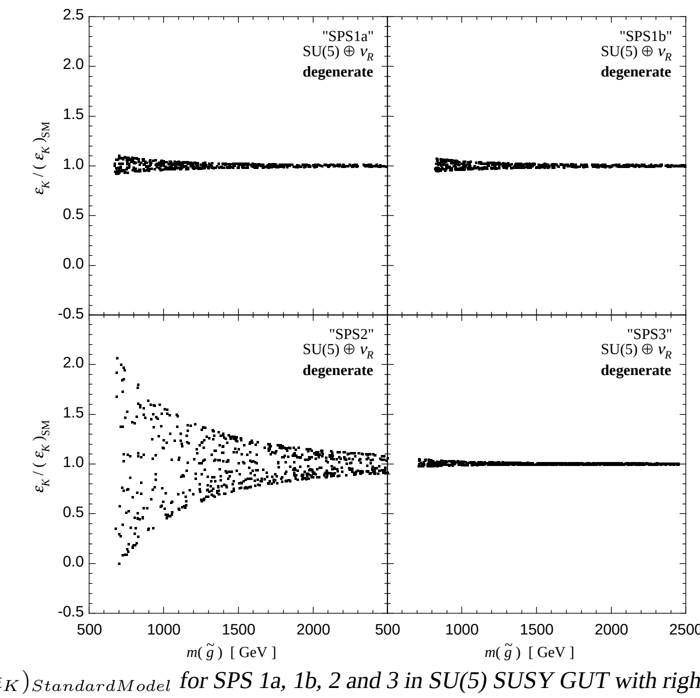

Figure 5-32. $\epsilon_K/(\epsilon_K)_{StandardModel}$ *for SPS 1a, 1b, 2 and 3 in SU(5) SUSY GUT with right-handed neutrinos for the degenerate right-handed neutrino mass case. Deviation is* $O(1)$ *only for SPS 2.*

Introduction

The Belle experiment in the KEK B Factory reported recently that the CP asymmetry in $B \to \phi K_S^0$ ($S_{\phi K_S^0}$) is $-0.96 \pm 0.50^{+0.09}_{-0.11}$, and 3.5σ deviation from the Standard-Model prediction 0.731 ± 0.056 is found [170]. At present the *BABAR* experiment does not observe such a large deviation, finding $0.45 \pm 0.43 \pm 0.07$ [171]. The combined result is not yet significant, however, Belle's result might be a signature of New Physics.

The CP violation in $B \to \phi K_S^0$ is sensitive to New Physics, since $b \to s\bar{s}s$ is a radiative process [172]. In fact, the SUSY models may predict a sizable deviation of the CP violation in $B \to \phi K_S^0$ from the Standard Model prediction. If the right-handed bottom and strange squarks have a sizable mixing, the gluon-penguin diagram may give a non-negligible contribution to $b \to s\bar{s}s$ in a broad parameter space where the contribution to $b \to s\gamma$ is a sub-dominant. $B \to \phi K_S^0$ in SUSY models has been studied in many papers [173][174][175][167].

In this article the correlation between the CP asymmetry in $B \to \phi K_S^0$ ($S_{\phi K_S^0}$) and the chromoelectric dipole moment (CEDM) of the strange quark (d_s^C) is discussed in SUSY models with right-handed squark mixing. In typical SUSY models, the left-handed squarks also have flavor mixing, due to the top-quark Yukawa coupling and the CKM mixing, and the left-handed bottom and strange squark mixing is as large as $\lambda^2 \sim 0.04$. When both the right-handed and left-handed squark mixings between the second and third generations are non-vanishing, a CEDM of the strange quark is generated. Since $S_{\phi K_S^0}$ and d_s^C may have a strong correlation in the SUSY models with the right-handed squark mixing, the constraint on d_s^C by the measurement of the EDM of ^{199}Hg limits the gluon-penguin contribution from the right-handed squark mixing to $S_{\phi K_S^0}$ [167].

In next section we discuss the ^{199}Hg EDM in SUSY models. In Section 5.3.3 the correlation between the CP asymmetry in $B \to \phi K_S^0$ and the CEDM of strange quark in the SUSY models with the right-handed squark mixing is presented. Section 5.3.3 is devoted to discussion.

The ^{199}Hg EDM in SUSY models

The EDMs of electron, neutron and nuclei are extensively studied in the SUSY models, and it is found that the relative phases among the flavor-diagonal A terms, B term in the Higgs potential, and the gaugino mass terms should be suppressed. However, even in that case, the EDMs are generated if both the left- and right-handed sfermions are mixed. It is especially interesting that EDM's are enhanced by heavier fermion masses, while they are suppressed by the mixing angles. Thus, the EDMs provide a stringent constraint on the SUSY models with both left- and right-handed sfermion mixings.

Before deriving the constraint on the bottom and strange squark mixing, we discuss the EDM of the nuclei. The EDMs of the diamagnetic atoms, such as ^{199}Hg, come from the CP-violating nuclear force by pion or eta meson exchange. The quark CEDMs,

$$H = \sum_{q=u,d,s} d_q^C \frac{i}{2} g_s \bar{q} \sigma^{\mu\nu} T^A \gamma_5 q G^A_{\mu\nu}, \tag{5.123}$$

generate the CP-violating meson-nucleon coupling, and the EDM of ^{199}Hg is evaluated in Ref. [176] as

$$d_{\rm Hg} = -3.2 \times 10^{-2} e \times (d_d^C - d_u^C - 0.012 d_s^C). \tag{5.124}$$

Chiral perturbation theory implies that $\bar{s}s$ in the matrix element of the nucleon is not suppressed, leading to a non-vanishing contribution from the CEDM of the strange quark. The suppression factor in front of d_s^C in Eq. (5.124) comes from the η meson mass and the CP-conserving coupling of the η meson and nucleon. From the current experimental bound on $d_{\rm Hg}$ ($d_{\rm Hg} < 2.1 \times 10^{-28} e\,{\rm cm}$) [177]:

$$e|d_d^C - d_u^C - 0.012 d_s^C| < 7 \times 10^{-27} e\,{\rm cm}\,. \tag{5.125}$$

If d_d^C and d_u^C are negligible in the equation,

$$e|d_s^C| < 7 \times 5.8 \times 10^{-25} e \,\mathrm{cm}\,. \tag{5.126}$$

The neutron EDM should also be affected by the CEDM of the strange quark. However, it is argued in Ref. [178] that this is suppressed by Peccei-Quinn symmetry. It is not clear at present whether the contribution of the CEDM of the strange quark is completely decoupled from the neutron EDM under Peccei-Quinn symmetry. In the following, we adopt the constraint on CEDM of the strange quark from ^{199}Hg.

In SUSY models, when the left-handed and right-handed squarks have mixings between the second and third generations, the CEDM of the strange quark is generated by a diagram in Fig. 5-33(a), and is enhanced by m_b/m_s. Using the mass insertion technique, d_s^C is given as

$$d_s^C = \frac{\alpha_s}{4\pi} \frac{m_{\tilde{g}}}{m_{\tilde{q}}^2} \left(-\frac{1}{3} N_1(x) - 3N_2(x)\right) \mathrm{Im}\left[(\delta_{LL}^{(d)})_{23}\,(\delta_{LR}^{(d)})_{33}\,(\delta_{RR}^{(d)})_{32}\right]\,, \tag{5.127}$$

up to the QCD correction, where $m_{\tilde{g}}$ and $m_{\tilde{q}}$ are the gluino and averaged squark masses. The functions N_i are given as

$$N_1(x) = \frac{3 + 44x - 36x^2 - 12x^3 + x^4 + 12x(2+3x)\log x}{(x-1)^6}, \tag{5.128}$$

$$N_2(x) = -2\frac{10 + 9x - 18x^2 - x^3 + 3(1 + 6x + 3x^2)\log x}{(x-1)^6}. \tag{5.129}$$

The mass insertion parameters $(\delta_{LL}^{(d)})_{23}$, $(\delta_{RR}^{(d)})_{32}$, and $(\delta_{LR}^{(d)})_{33}$ are given by

$$(\delta_{LL}^{(d)})_{23} = \frac{\left(m^2_{\tilde{d}_L}\right)_{23}}{m_{\tilde{q}}^2},\ (\delta_{RR}^{(d)})_{32} = \frac{\left(m^2_{\tilde{d}_R}\right)_{32}}{m_{\tilde{q}}^2},\ (\delta_{LR}^{(d)})_{33} = \frac{m_b\,(A_b - \mu\tan\beta)}{m_{\tilde{q}}^2}, \tag{5.130}$$

where $(m^2_{\tilde{d}_{L(R)}})$ is the left-handed (right-handed) down-type squark mass matrix. In typical SUSY models, $(\delta_{LL}^{(d)})_{23}$ is $O(\lambda^2) \simeq 0.04$. From this formula, d_s^C is estimated in a limit of $x \to 1$ as

$$ed_s^C = e\frac{\alpha_s}{4\pi} \frac{m_{\tilde{g}}}{m_{\tilde{q}}^2} \left(-\frac{11}{30}\right) \mathrm{Im}\left[(\delta_{LL}^{(d)})_{23}\,(\delta_{LR}^{(d)})_{33}\,(\delta_{RR}^{(d)})_{32}\right] \tag{5.131}$$

$$= -4.0 \times 10^{-23} \sin\theta\, e\,\mathrm{cm} \left(\frac{m_{\tilde{q}}}{500\mathrm{GeV}}\right)^{-3} \left(\frac{(\delta_{LL}^{(d)})_{23}}{0.04}\right) \left(\frac{(\delta_{RR}^{(d)})_{32}}{0.04}\right) \left(\frac{\mu\tan\beta}{5000\mathrm{GeV}}\right)\,, \tag{5.132}$$

where $\theta = \arg[(\delta_{LL}^{(d)})_{23}\,(\delta_{LR}^{(d)})_{33}\,(\delta_{RR}^{(d)})_{32}]$. Here, we neglect the contribution proportional to A_b, since it is subdominant. From this formula, it is obvious that the right-handed squark mixing or the CP-violating phase should be suppressed. For example, for $m_{\tilde{q}} = 500$GeV, $\mu\tan\beta = 5000$GeV, and $(\delta_{LL}^{(d)})_{23} = 0.04$,

$$|\sin\theta(\delta_{RR}^{(d)})_{32}| < 5.8 \times 10^{-4}. \tag{5.133}$$

Correlation between d_s^C and $B \to \phi K_S^0$ in models with right-handed squark mixing

Let us discuss the correlation between d_s^C and $S_{\phi K_S^0}$ in the SUSY models with right-handed squark mixing. As mentioned in Introduction, the right-handed bottom and strange squark mixing may lead to the sizable deviation of $S_{\phi K_S^0}$ from the Standard Model prediction by the gluon-penguin diagram, especially for large $\tan\beta$. The box diagrams with the right-handed squark mixing also contribute to $S_{\phi K_S^0}$, but they tend to be sub-dominant, and do not generate a large deviation of $S_{\phi K_S^0}$ from the Standard Model prediction. Thus, for simplicity, we will neglect the box contribution in this article.

The effective operator inducing the gluon-penguin diagram by the right-handed squark mixing is

$$H = -C_8^R \frac{g_s}{8\pi^2} m_b \overline{s_R} \sigma^{\mu\nu} T^A b_L G^A_{\mu\nu}. \tag{5.134}$$

When the right-handed squarks are mixed, the dominant contribution to C_8^R is supplied by a diagram with the double mass insertion of $(\delta^{(d)}_{RR})_{32}$ and $(\delta^{(d)}_{RL})_{33}$ (Fig. 5-34(b)). This contribution is specially significant when $\mu\tan\beta$ is large. The contribution of Fig. 5-33(b) to C_8^R is given as

$$C_8^R = \frac{\pi\alpha_s}{m_{\tilde{q}}^2}\frac{m_{\tilde{g}}}{m_b}(\delta^{(d)}_{LR})_{33}(\delta^{(d)}_{RR})_{32}(-\frac{1}{3}M_1(x) - 3M_2(x)) \tag{5.135}$$

up to QCD corrections. Here,

$$M_1(x) = \frac{1 + 9x - 9x^2 - x^3 + (6x + 6x^2)\log x}{(x-1)^5}, \tag{5.136}$$

$$M_2(x) = -2\frac{3 - 3x^2 + (1 + 4x + x^2)\log x}{(x-1)^5}. \tag{5.137}$$

In a limit of $x \to 1$, C_8^R is reduced to

$$C_8^R = \frac{7\pi\alpha_s}{30 m_b m_{\tilde{q}}}(\delta^{(d)}_{LR})_{33}(\delta^{(d)}_{RR})_{32}. \tag{5.138}$$

Comparing Eq. (5.131) and Eq. (5.138), we see a strong correlation between d_s^C and C_8^R:

$$d_s^C = -\frac{m_b}{4\pi^2}\frac{11}{7}\mathrm{Im}\left[(\delta^{(d)}_{LL})_{23} C_8^R\right], \tag{5.139}$$

up to QCD corrections. The coefficient $11/7$ in Eq. (5.139) changes from 3 to 1 for $0 < x < \infty$.

In Fig. 5-34 the correlation between d_s^C and $S_{\phi K_S^0}$ is presented. Here, we assume $d_s^C = -m_b/(4\pi^2)\mathrm{Im}[(\delta^{(d)}_{LL})_{23}C_8^R]$, up to QCD corrections. Here, we take $(\delta^{(d)}_{LL})_{23} = -0.04$, $\arg[C_8^R] = \pi/2$ and $|C_8^R|$ corresponding to $10^{-5} < |(\delta^{(d)}_{RR})_{32}| < 0.5$. The matrix element of chromomagnetic moment in $B \to \phi K_S^0$ is

$$\langle \phi K_S^0 | \frac{g_s}{8\pi^2} m_b (\bar{s}_i \sigma^{\mu\nu} T^a_{ij} P_R b_j) G^a_{\mu\nu} | \overline{B}_d \rangle = \kappa \frac{4\alpha_s}{9\pi}(\epsilon_\phi p_B) f_\phi m_\phi^2 F_+(m_\phi^2), \tag{5.140}$$

and $\kappa = -1.1$ in the heavy-quark effective theory [175]. Since κ may suffer from large hadronic uncertainties, we take $\kappa = -1$ and -2. From this figure, it is found that the deviation of $S_{\phi K_S^0}$ from the Standard Model prediction due to the gluon penguin contribution should be tiny when the constraint on d_s^C in Eq. (5.126) is applied.

Discussion

In this article the correlation between the CP asymmetry in $B \to \phi K_S^0$ and the chromoelectric dipole moment (CEDM) of strange quark has been discussed in SUSY models with right-handed squark mixing. While the gluon-penguin diagram might give a large deviation of $S_{\phi K_S^0}$ from the Standard Model prediction, the size is limited from the constraint on the CEDM of the strange quark. The constraint from the CEDM of the strange quark on the mixing between the right-handed strange and bottom squark is the most stringent at present, compared with other processes where the left-handed squarks are also mixed. For example, the $\mathcal{B}b \to s\gamma$ gives the constraint as

$$\left|(\delta^{(d)}_{RR})_{23}\right| \lesssim 0.27 \left(\frac{\mu\tan\beta}{5000GeV}\right)^{-1}\left(\frac{m_{\tilde{q}}}{500GeV}\right)^2, \tag{5.141}$$

which is looser. Also, the right-handed down-type squark mixing is related to the left-handed slepton mixing in the SUSY SU(5) GUT, and the experimental bound on $\mathcal{B}(\tau \to \mu\gamma)$ gives a constraint on the mixing between the right-handed strange and bottom squark [174]. While the current bound on $\mathcal{B}(\tau \to \mu\gamma)$ may exclude the possibility of a large deviation of $S_{\phi K_S^0}$, a sizable deviation is still allowed.

It has been argued recently in Ref. [179] that the measurement of the deuteron EDM may improve the bound on the CP-violating nuclear force by two orders of magnitude. If this is realized, it will be a stringent test of SUSY models with right-handed squark mixing, such as SUSY GUTs.

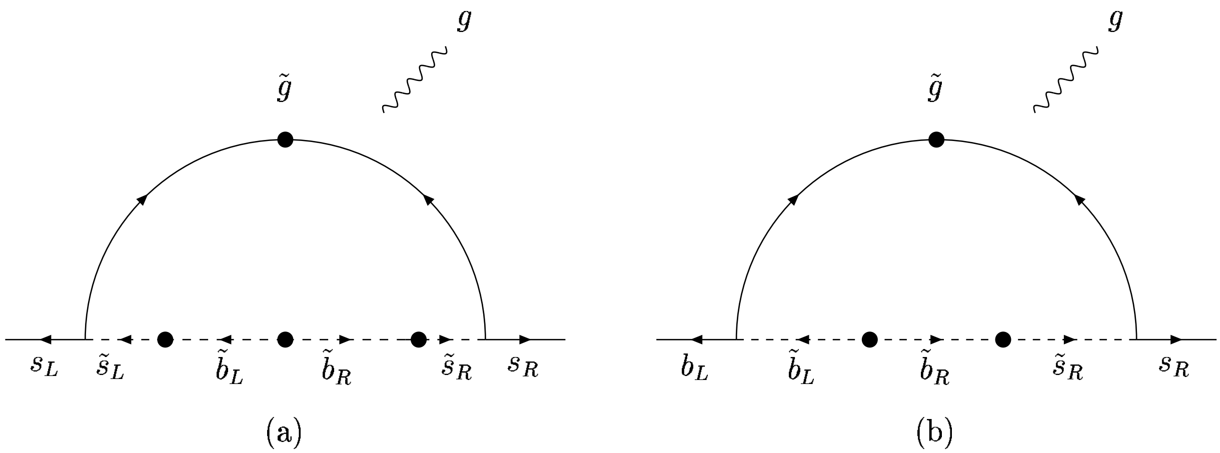

Figure 5-33. *a) The dominant diagram contributing to the CEDM of the strange quark when both the left-handed and right-handed squarks have mixings. b) The dominant SUSY diagram contributing to the CP asymmetry in $B \to \phi K_S^0$ when the right-handed squarks have a mixing.*

5.3.4 SUSY Analysis in B Decays: the Mass Insertion Approximation

>– M. Ciuchini, A. Masiero, L. Silvestrini, S. K. Vempati and O. Vives –<

Introduction

Our knowledge of the flavor sector in the Minimal Supersymmetric extension of the Standard Model (MSSM) is still very limited. Only after the discovery of SUSY particles and the measurement of the supersymmetric spectrum we will be able to explore in detail this fundamental piece of the MSSM. Nevertheless, we already have a lot of useful information on this sector from experiments looking for indirect effects of SUSY particles in low-energy experiments [180, 181].

To analyze flavor-violating constraints at the electroweak scale, the model-independent mass-insertion (MI) approximation is advantageous [182, 183, 184, 185]. In this method, the experimental limits lead to upper bounds on the parameters (or combinations of) $\delta^f_{ij} \equiv \Delta^f_{ij}/m^2_{\tilde{f}}$; where Δ^f_{ij} is the flavor-violating off-diagonal entry appearing in the $f = (u, d, l)$ sfermion mass matrices in the basis of diagonal Yukawa matrices and $m^2_{\tilde{f}}$ is the average sfermion mass. In addition, the mass-insertions are further sub-divided into $LL/LR/RL/RR$ types, labeled by the chirality of the corresponding Standard Model fermions. With the help of this MI formalism we can easily estimate the sensitivity of different processes to offdiagonal entries in the sfermion mass matrices. In this respect, it is instructive to compare the sensitivity of kaon and B physics experiments.

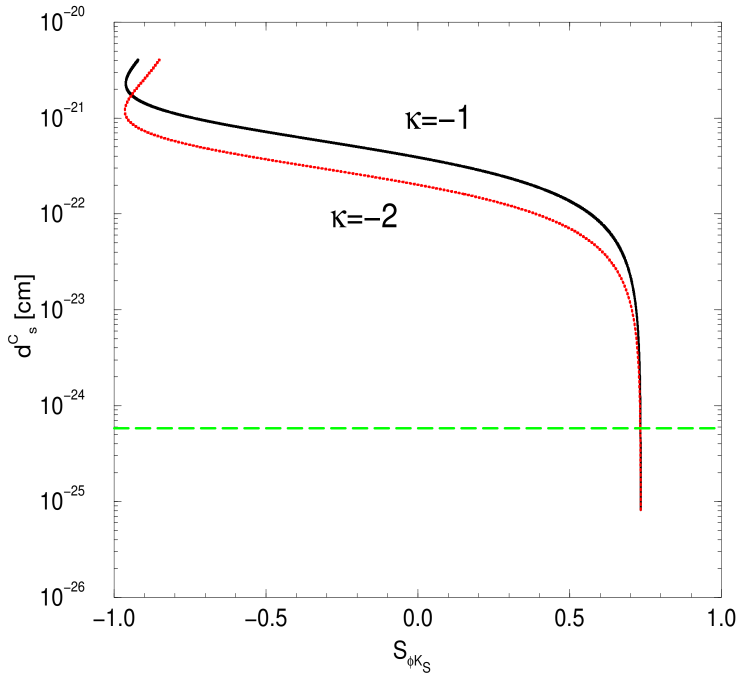

Figure 5-34. *The correlation between d^C_s and $S_{\phi K^0_S}$ assuming $d^C_s = -m_b/(4\pi^2)\text{Im}[(\delta^{(d)}_{LL})_{23} C^R_8]$. Here, $(\delta^{(d)}_{LL})_{23} = -0.04$ and $\arg[C^R_8] = \pi/2$. κ comes from the matrix element of chromomagnetic moment in $B \to \phi K^0_S$. The dashed line is the upperbound on d^C_s from the EDM of ^{199}Hg atom.*

If we assume that indirect CP violation in the kaon sector gets a sizable contribution from SUSY, while the kaon mass difference is mainly due to the Standard Model loops we have,

$$2.3 \times 10^{-3} \geq \varepsilon_K^{\text{SUSY}} = \frac{\text{Im}\, M_{12}|_{\text{SUSY}}}{\sqrt{2}\,\Delta M_K|_{\text{SM}}} \simeq \frac{\alpha_s^2}{\alpha_W^2}\,\frac{M_W^2}{M^2_{\text{SUSY}}}\,\frac{\text{Im}\left\{(\delta^d_{12})^2_{LL}\right\}}{(V_{cd}V^*_{cs})^2\,\frac{m_c^2}{M_W^2}}$$

$$\simeq 12.5 \times 0.026 \times \frac{\text{Im}\left\{(\delta^d_{12})^2_{LL}\right\}}{1.5 \times 10^{-5}} \Rightarrow \sqrt{\text{Im}\left\{(\delta^d_{12})^2_{LL}\right\}} \leq 3.3 \times 10^{-4}\,, \tag{5.142}$$

where we have assumed a SUSY mass scale of 500 GeV. In the Standard Model contribution, we have taken into account that quark masses must be present because of the GIM mechanism, and we have used the fact that the loop function in the W diagram, $S(x_c = m_c^2/M_W^2) \simeq x_c$ for $x_c \ll 1$ [186][3]. We have ignored factors of $\mathcal{O}(1)$, as well as the different loop functions in the gluino contributions. Hence, we can see that the SUSY contribution is suppressed by the heavy squark masses with respect to the W boson mass. However, the SUSY gluino contribution is proportional to the strong coupling, while the Standard Model contribution is proportional to the weak coupling. Apart from these factors, we have to compare the mass insertion $(\delta^d_L)_{12}$ with $V_{cd}\frac{m_c}{M_W}V^*_{cs}$. Hence we can see that, due to the small fermion masses and/or mixing angles, ε_K is in fact sensitive to MI at the level of a few $\times\, 10^{-4}$ [187].

[3]Here the charm–W loop gives the main contribution for the kaon mass difference and this is sufficient for our estimate. Long distance (and top quark) effects are not included here, although they give a sizable ($\simeq 30\%$) contribution to Δm_K

Table 5-14. *Bounds on the mass insertions from ε_K, ε'/ε, BR($b \to s\gamma$) and $\Delta M_{B_{d,s}}$ for $m_{\tilde{q}} = 500$ GeV. For different squark masses, bounds on $(\delta_{LR})_{12}$ and $(\delta_{LR})_{13}$ scale as $(m_{\tilde{q}}(\mathrm{GeV})/500)^2$, while other bounds scale as $(m_{\tilde{q}}(\mathrm{GeV})/500)$. These bounds are equal under the exchange $L \leftrightarrow R$.*

x	$\sqrt{\left\lvert \mathrm{Im}\left(\delta^d_{12}\right)^2_{LL}\right\rvert}$	$\sqrt{\left\lvert \mathrm{Im}\left(\delta^d_{12}\right)^2_{LR}\right\rvert}$	$\sqrt{\left\lvert \mathrm{Im}\left(\delta^d_{12}\right)_{LL}\left(\delta^d_{12}\right)_{RR}\right\rvert}$	$\sqrt{\left\lvert \mathrm{Re}\left(\delta^d_{13}\right)^2_{LL}\right\rvert}$	$\sqrt{\left\lvert \mathrm{Re}\left(\delta^d_{13}\right)^2_{LR}\right\rvert}$
0.3	2.9×10^{-3}	1.1×10^{-5}	1.1×10^{-4}	4.6×10^{-2}	5.6×10^{-2}
1.0	6.1×10^{-3}	2.0×10^{-5}	1.3×10^{-4}	9.8×10^{-2}	3.3×10^{-2}
4.0	1.4×10^{-2}	6.3×10^{-5}	1.8×10^{-4}	2.3×10^{-1}	3.6×10^{-2}
	$\sqrt{\left\lvert \mathrm{Re}\left(\delta^d_{23}\right)^2_{LL}\right\rvert}$	$\left\lvert\left(\delta^d_{23}\right)_{LR}\right\rvert$	$\sqrt{\left\lvert \mathrm{Re}\left(\delta^d_{23}\right)_{LL}\left(\delta^d_{23}\right)_{RR}\right\rvert}$	$\sqrt{\left\lvert \mathrm{Re}\left(\delta^d_{13}\right)_{LL}\left(\delta^d_{13}\right)_{RR}\right\rvert}$	
0.3	0.21	1.3×10^{-2}	7.4×10^{-2}	1.6×10^{-2}	
1.0	0.45	1.6×10^{-2}	8.3×10^{-2}	1.8×10^{-2}	
4.0	1	3.0×10^{-2}	1.2×10^{-1}	2.5×10^{-2}	

Similarly, we analyse the MI in (1,3) transitions from the B^0 CP asymmetries,

$$0.74 \geq a_{J/\psi}\big|_{\text{SUSY}} = \frac{\mathrm{Im}\, M_{12}|_{\text{SUSY}}}{|M_{12}|_{\text{SM}}} \simeq \frac{\alpha_s^2}{\alpha_W^2}\,\frac{M_W^2}{M_{\text{SUSY}}^2}\,\frac{\mathrm{Im}\left\{(\delta^d_{13})^2_{LL}\right\}}{(V_{tb}V^*_{td})^2\,\frac{m_t^2}{M_W^2}}$$

$$\simeq 12.5 \times 0.026 \times \frac{\mathrm{Im}\left\{(\delta^d_{13})^2_{LL}\right\}}{3 \times 10^{-4}} \Rightarrow \sqrt{\mathrm{Im}\left\{(\delta^d_{13})^2_{LL}\right\}} \leq 0.0\,,3 \qquad (5.143)$$

where again we use $S(x_t) \simeq x_t$ [186]. In this case we have some differences with respect to the situation in the kaon sector. First, in this case the combination of fermion masses and mixing angles in the Standard Model contribution is larger by a factor of 20. So, if we could reach the same experimental sensitivity in B CP experiments as in kaon CP violation experiments we would be able to explore MI a factor $\sqrt{20}$ larger. However, the main difference between both experiments is the different experimental sensitivity to the observables. In kaon physics we are sensitive to signals of CP violation that are 10^3 times smaller than the kaon mass difference. In B physics we measure CP violation effects of the same order as the B mass difference, and we are sensitive to signals roughly one order of magnitude smaller. This is the main reason why CP experiments in kaon physics are sensitive to much smaller entries in the sfermion mass matrices than experiments in B physics [187]. We can compare these estimates with the actual bounds in Table 5-14 and we see that our simplified calculations are correct as order of magnitude estimates.

However, this does not mean that it is impossible to find signs of supersymmetry in B physics experiments. We have several reasons to expect larger off diagonal entries in the elements associated with $b \to s$ or $b \to d$ transitions than in $s \to d$ transitions. In some grand unified models, large atmospheric neutrino mixing is associated with large right-handed down quark mixing [188, 189, 190]. Sizable mixing in transitions between the third and second generations is also generically expected in flavor models [191]. In fact, we have only weak experimental constraints on $b \to s$ transitions from the $\mathcal{B}(b \to s\gamma$ and Δm_s, as shown in Table 5-14. So, the question is now, are large SUSY effects possible in B transitions?.

FCNC in GUT supersymmetry

In a SUSY GUT, quarks and leptons are in the same multiplet. As long as the scale associated with the transmission of SUSY breaking to the visible sector is larger than the GUT breaking scale, the quark-lepton unification also seeps into the SUSY breaking soft sector, leading to squark-slepton mass-squared unification [192]. The exact relations between the mass matrices depend on the choice of the GUT gauge group. For instance, in SU(5) $(\Delta^d_{ij})_{RR}$ and $(\Delta^l_{ij})_{LL}$ are equal; in SO(10) all Δ_{ij} are equal at M_{GUT} implying strong correlations within FCNCs at that scale that can have significant implications on flavor phenomenology.

To be specific, we concentrate on SUSY SU(5), with soft terms generated above M_{GUT}. We assume generic flavor-violating entries to be present in the sfermion matrices at the GUT scale [4]. The part of the superpotential relevant for quarks and charged lepton masses can be written as

$$W_{\rm SU(5)} = h^u_{ij}\, T_i\, T_j\, H + h^d_{ij} T_i\, \overline{F}_j\, \overline{H} + \mu\, H\, \overline{H}, \tag{5.144}$$

where we have used the standard notation, with T transforming as a 10 and $\overline{F}$ as a $\overline{5}$ under SU(5). The corresponding SU(5) invariant soft potential has now the form:

$$-\mathcal{L}_{soft} = m^2_{T_{ij}} \tilde{T}^\dagger_i \tilde{T}_j + m^2_{\overline{F}} \tilde{\overline{F}}^\dagger_i \tilde{\overline{F}}_j + m^2_H H^\dagger H + m^2_{\overline{H}} \overline{H}^\dagger \overline{H} + A^u_{ij} T_i T_j H + A^d_{ij} T_i \overline{F}_j \overline{H} + B\mu H \overline{H}\,. \tag{5.145}$$

Rewriting this in terms of the Standard Model representations, we have

$$-\mathcal{L}_{soft} = m^2_{Q_{ij}} \tilde{Q}^\dagger_i \tilde{Q}_j + m^2_{u^c_{ij}} \tilde{u^c}^\star_i \tilde{u^c}_j + m^2_{e^c_{ij}} \tilde{e^c}^\star_i \tilde{e^c}_j + m^2_{d^c_{ij}} \tilde{d^c}^\star_i \tilde{d^c}_j + m^2_{L_{ij}} \tilde{L}^\dagger_i \tilde{L}_j +$$
$$m^2_{H_1} H^\dagger_1 H_1 + m^2_{H_2} H^\dagger_2 H_2 + A^u_{ij}\, \tilde{Q}_i \tilde{u^c}_j H_2 + A^d_{ij}\, \tilde{Q}_i \tilde{d^c}_j H_1 + A^e_{ij}\, \tilde{L}_i \tilde{e^c}_j H_1 + \ldots \tag{5.146}$$

$$m^2_Q = m^2_{\tilde{e^c}} = m^2_{\tilde{u^c}} = m^2_T, \qquad m^2_{\tilde{d^c}} = m^2_L = m^2_{\overline{F}}, \qquad A^e_{ij} = A^d_{ji}\,. \tag{5.147}$$

Eqs. (5.147) are matrices in flavor space. These equations lead to relations within the slepton and squark flavor-violating off-diagonal entries Δ_{ij}. These are:

$$(\Delta^u_{ij})_L = (\Delta^u_{ij})_R = (\Delta^d_{ij})_L = (\Delta^l_{ij})_R, \quad (\Delta^d_{ij})_R = (\Delta^l_{ij})_L, \quad (\Delta^d_{ij})_{LR} = (\Delta^l_{ji})_{LR} = (\Delta^l_{ij})^\star_{RL}\,. \tag{5.148}$$

These relations are exact at M_{GUT}; however, after SU(5) breaking, quarks and leptons suffer different renormalization effects and are thus altered at M_W. It is easy to see from the RG equations that off-diagonal elements in the squark mass matrices in the first two of Eqs. (5.148) are approximately not renormalized due to the smallness of CKM mixing angles and that the sleptonic entries in them are left unchanged (in the absence of right-handed neutrinos). On the other hand, the last equation receives corrections due to the different nature of the RG scaling of the LR term (A-parameter). This correction can be roughly approximated as proportional to the corresponding fermion masses. Taking this into consideration, we can now rewrite the Eqs. (5.148) at the weak scale,

$$(\delta^d_{ij})_{RR} \;\approx\; \frac{m^2_L}{m^2_{d^c}} (\delta^l_{ij})_{LL}, \qquad (\delta^{u,d}_{ij})_{LL} \;\approx\; \frac{m^2_{e^c}}{m^2_Q} (\delta^l_{ij})_{RR},$$
$$(\delta^u_{ij})_{RR} \;\approx\; \frac{m^2_{e^c}}{m^2_{u^c}} (\delta^l_{ij})_{RR}, \qquad (\delta^d_{ij})_{LR} \;\approx\; \frac{m^2_{L_{avg}}}{m^2_{Q_{avg}}}\, \frac{m_b}{m_\tau}\, (\delta^l_{ij})^\star_{RL}, \tag{5.149}$$

where $m^2_{L_{avg}}$ ($m^2_{Q_{avg}}$) are given by the geometric average of left- and right-handed slepton (down-squark) masses $\sqrt{m^2_L\, m^2_{e^c}}$ $\left(\sqrt{m^2_Q\, m^2_{d^c}}\right)$, all defined at the weak scale.

To account for neutrino masses, we can use the seesaw mechanism by adding singlet right-handed neutrinos. In their presence, additional couplings occur in Eqs. (5.144 - 5.148) at the high scale which affect the RG evolution of slepton matrices. To understand the effect of these new couplings, one can envisage two scenarios [193]: (a) small couplings and/or small mixing in the neutrino Dirac Yukawa matrix, (b) large couplings and large mixing in the neutrino sector. In case (a), the effect on slepton mass matrices due to neutrino Dirac Yukawa couplings is very small and the above relations Eqs. (5.149) still hold. In case (b), however, large RG effects can significantly modify the slepton doublet flavor structure while keeping the squark sector and right handed charged slepton matrices essentially unmodified, thus breaking the GUT-symmetric relations. Even in this case, barring accidental cancellations among the mass insertions already present at M_{GUT} and the radiatively generated mass insertions between M_{GUT} and M_{ν_R}, there exists an upper bound on the down quark δ parameters of the form:

$$|(\delta^d_{ij})_{RR}| \quad \leq \quad \frac{m^2_L}{m^2_{d^c}} |(\delta^l_{ij})_{LL}|\,, \tag{5.150}$$

[4]Note that even assuming complete universality of the soft breaking terms at M_{Planck}, as in mSUGRA, the RG effects on M_{GUT} will induce flavor off-diagonal entries at the GUT scale [194].

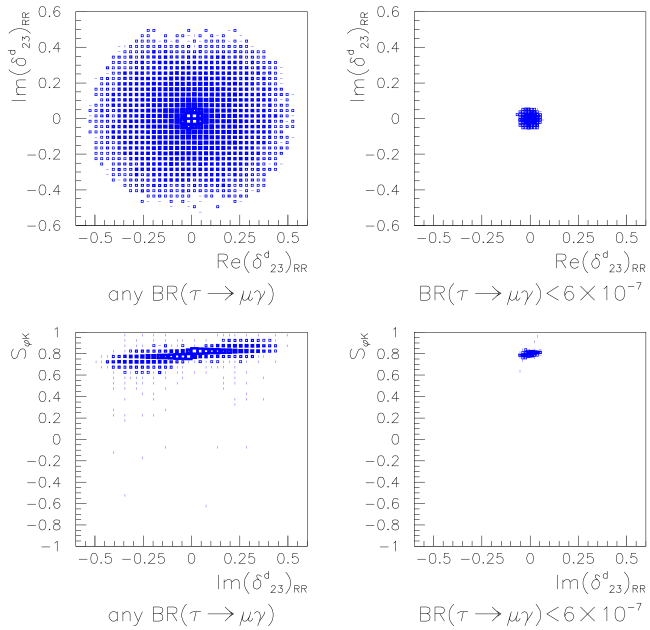

Figure 5-35. *Allowed regions in the Re$(\delta^d_{23})_{RR}$–Im$(\delta^d_{23})_{RR}$ plane and in the $S_{K\phi}$–Im$(\delta^d_{23})_{RR}$ plane. Constraints from $B \to X_s\gamma$, $BR(B \to X_s\ell^+\ell^-)$, and the lower bound on Δm_s have been used.*

while the last three equations in Eq. (5.149) remain valid in this case.

The relations (5.149, 5.150) predict links between lepton and quark flavor-changing transitions at the weak scale. For example, we see that $\mu \to e\gamma$ can be related to $K^0 - \overline{K}^0$ mixing and to $D^0 - \overline{D}^0$ mixing. Similarly, one can expect correlations between $\tau \to e\gamma$ and $B_d - \overline{B}_d$ mixing, as well as between $\tau \to \mu\gamma$ and $b \to s$ transitions such as $B \to \phi K^0_S$.

To show the impact of these relations, let us assume that all the flavor diagonal sfermion masses are approximately universal at the GUT scale, with $m^2_T = m^2_{\tilde{F}} = m^2_H = m^2_{\tilde{H}} = m^2_0$, with flavor off-diagonal entries $m^2_{\tilde{f}} = m^2_0 \mathbf{1} + \mathbf{\Delta^f_{ij}}$, with $|\Delta^f_{ij}| \le m^2_0$. Δ^f_{ij} can be present either through the running from the Planck scale to the GUT scale [194] or through some flavor non-universality originally present [195, 196, 197]. All gaugino masses are unified to $M_{1/2}$ at M_{GUT}. For a given set of initial conditions ($M_{1/2}$, m^2_0, A_0, Δ_{ij}, $\tan\beta$) we obtain the full spectrum at M_W with the requirement of radiative symmetry breaking. We then apply limits from direct searches on SUSY particles. Finally, we calculate the contributions of different δ_{23} parameters to both leptonic and hadronic processes, considering the region in the $(m_0, M_{1/2})$ plane corresponding to a relatively light sparticle spectrum, with squark masses of roughly 350–550 GeV and slepton masses of about 150–400 GeV.

$b \to s$ transitions have recently received much attention, as it has been shown that the discrepancy with Standard Model expectations in the measurements of $A_{CP}(B \to \phi K^0_S)$ can be attributed to the presence of large neutrino mixing within SO(10) models [188, 190, 189]. Subsequently, a detailed analysis has been presented [198, 199] within the context of the MSSM. It has been shown that, for squark and gluino masses around 350 GeV, the presence of a $\mathcal{O}(1)$ $(\delta^d_{23})_{LL,RR}$ could lead to significant discrepancies from the Standard Model expectations. Similar statements hold for a $\mathcal{O}(10^{-2})$ LR or RL MI. Here, we study the impact of LFV bounds on these δ parameters and subsequently the effect on B physics observables. In Table 5-15, we present upper bounds on δ^d_{23} within the above mass ranges for three values of the limits on $\mathcal{B}(\tau \to \mu\gamma)$. There are no bounds on $(\delta^d_{23})_{LL}$ because, as is well known [200, 201, 202],

Table 5-15. *Bounds on (δ^d_{23}) from $\mathcal{B}(\tau \to \mu\,\gamma)$ for three values of the branching ratios for $\tan\beta$ = 10.*

Type	$<\ 6 \cdot 10^{-7}$	$<\ 1 \cdot 10^{-7}$	$<\ 1 \cdot 10^{-8}$
LL	-	-	-
RR	0.070	0.030	0.010
RL	0.080	0.035	0.010
LR	0.080	0.035	0.010

large values of $(\delta^l_{ij})_{RR}$ are still allowed, due to the cancellations of bino and higgsino contributions in the decay amplitude.

At present, the constraints coming from B physics are stronger than those obtained for the lepton sector in the cases of $(\delta^d_{23})_{LL,LR,RL}$. Therefore no impact on B phenomenology is expected even if the present bound on $\mathcal{B}(\tau \to \mu\,\gamma)$ were pushed down to 1×10^{-7}. On the contrary, the bound on $(\delta^d_{23})_{RR}$ induced by $\mathcal{B}(\tau \to \mu\,\gamma)$ is already at present much stronger than the bounds from hadronic processes, reducing considerably the room left for SUSY effects in B decays. To illustrate this point in detail, we repeat the analysis of Ref. [199] including the bounds coming from lepton processes. We therefore compute at the NLO branching ratios and CP asymmetries for $B \to X_s\gamma$ and $B \to \phi K^0_s$, $\mathcal{B}(B \to X_s\ell^+\ell^-)$ and Δm_s (see Ref. [199] for details). In the first row of Fig. 5-35, we plot the probability density in the Re$(\delta^d_{23})_{RR}$–Im$(\delta^d_{23})_{RR}$ plane for different upper bounds on $\mathcal{B}(\tau \to \mu\,\gamma)$. Note that making use of Eq. (5.149) with $|(\delta^l_{23})_{LL}|\ <\ 1$, implies $|(\delta^d_{23})_{RR}| \lesssim 0.5$ as the ratio $(m^2_L/m^2_{d^c})$ varies roughly between $(0.2-0.5)$ at the weak scale, for the chosen high scale boundary conditions. The effect on $(\delta^d_{23})_{RR}$ of the upper bound on $\mathcal{B}(\tau \to \mu\,\gamma)$ is dramatic already with the present experimental value. Correspondingly, as can be seen from the second row of Fig. 5-35, the possibility of large deviations from the Standard Model in the coefficient $S_{\phi K}$ of the sine term in the time-dependent $A_{CP}(B \to \phi K^0_S)$ is excluded in the RR case. Hence, we conclude that in SUSY GUTs the most likely possibility to strongly depart from the Standard Model expectations for $S_{\phi K}$ relies on a sizable contribution from $(\delta^d_{23})_{LL}$ or $(\delta^d_{23})_{LR,RL}$, as long as they are small enough to be within the severe limits imposed by $\mathcal{B}(B \to X_s\gamma)$ [199]. These results would not change significantly if one started with a SO(10) theory instead of a SU(5) theory. The relation in Eq. (5.149) would be still valid however with the additional constraint: $(\delta^d_{ij})_{RR} = m^2_Q/m^2_d(\delta^d_{ij})_{LL} = m^2_L/m^2_d(\delta^l_{ij})_{LL}$. The results of our analysis are therefore valid also for SO(10), although stronger correlations are generally expected.

Exploiting the Grand Unified structure of the theory, we can obtain similar bounds on other δ^d_{ij} parameters. For example, considering the first two generations, the bound on δ^d_{12} from $\mathcal{B}(\mu \to e\,\gamma)$ can, in many cases, compete with the bound from ΔM_K [185]. Similar comparisons can be made for the δ^d_{13} from limits on $\mathcal{B}(\tau \to e\,\gamma)$ and $B^0_d-\overline{B}^0_d$ mixing.

Example: SU(3) flavor theory

Finally as an example, we discuss here the main features of a recent supersymmetric SU(3) flavor model [203, 204, 205] that successfully reproduces quark and lepton masses and mixing angles, and predicts the structure of the sfermion mass matrices. Under this SU(3) family symmetry, all left-handed fermions (ψ_i and ψ^c_i) are triplets. To allow for the spontaneous symmetry breaking of SU(3), it is necessary to add several new scalar fields that are either triplets ($\overline{\theta}_3$, $\overline{\theta}_{23}$, $\overline{\theta}_2$) or anti-triplets (θ_3, θ_{23}). We assume that SU(3)$_F$ is broken in two steps. The first step occurs when θ_3 gets a large vev breaking SU(3) to SU(2). Subsequently, a smaller vev of θ_{23} breaks the remaining symmetry. After this breaking, we obtain the effective Yukawa couplings at low energies through the Froggatt-Nielsen mechanism. In this theory, the Yukawa superpotential is

$$W_Y = H\psi_i\psi^c_j\left[\theta^i_3\theta^j_3\ +\ \theta^i_{23}\theta^j_{23}\Sigma\ +\ \left(\epsilon^{ikl}\overline{\theta}_{23,k}\overline{\theta}_{3,l}\theta^j_{23}\left(\theta_{23}\overline{\theta_3}\right)\ +\ (i \leftrightarrow j)\right)\right]\ , \tag{5.151}$$

and so the Yukawa textures are

$$Y^f \propto \begin{pmatrix} 0 & \varepsilon^3 e^{i\delta} X_1 & \varepsilon^3 e^{i(\delta+\beta_3)} X_2 \\ \dots & \varepsilon^2 \frac{\Sigma}{|a_3|^2} & \varepsilon^2 e^{i\beta_3} \frac{\Sigma}{|a_3|^2} \\ \dots & \dots & e^{2i\chi} \end{pmatrix}, \tag{5.152}$$

where $\overline{\varepsilon} = \langle\theta_{23}\rangle/M \simeq 0.15$ with M a mediator mass in terms of dimension greater than three, and the X_a are $\mathcal{O}(1)$ coefficients. In the same way, the supergravity Kähler potential receives new contributions after SU(3)$_F$ breaking,

$$K = \psi_i^\dagger \psi_j \Big(\delta^{ij}(c_0 + d_0 XX^\dagger) + \frac{1}{M^2}[\theta_3^{i\dagger}\theta_3^j(c_1 + d_1 XX^\dagger) + \theta_{23}^{i\dagger}\theta_{23}^j(c_2 + d_2 XX^\dagger)] + \\ +(\epsilon^{ikl}\overline{\theta}_{3,k}\overline{\theta}_{23,l})^\dagger(\epsilon^{jmn}\overline{\theta}_{3,m}\overline{\theta}_{23,n})(c_3 + d_3 XX^\dagger)]\Big) , \tag{5.153}$$

where c_i, d_i are $\mathcal{O}(1)$ coefficients and we include a field X with non-vanishing F-term. From here we obtain the structure of the sfermion mass matrices [203, 204, 205]. In the basis of diagonal quark Yukawa couplings (SCKM basis) we obtain for the down quarks, suppressing factors $\mathcal{O}(1)$,

$$(M^2_{\tilde{D}_R})^{\rm SCKM} \simeq \begin{pmatrix} 1+\overline{\varepsilon}^3 & -\overline{\varepsilon}^3 e^{-i\omega} & -\overline{\varepsilon}^3 e^{-i\omega} \\ -\overline{\varepsilon}^3 e^{i\omega} & 1+\overline{\varepsilon}^2 & \overline{\varepsilon}^2 \\ -\overline{\varepsilon}^3 e^{i\omega} & \overline{\varepsilon}^2 & 1+\overline{\varepsilon} \end{pmatrix} m_0^2 \, . \tag{5.154}$$

Thus, we can see that in $3 \to 2$ transitions we have off-diagonal entries of order $\overline{\varepsilon}^2$, although these must still be small to have large effects. In the case of lepton flavor violation, the slepton mass matrices are similar to Eq. (5.154) with different $\mathcal{O}(1)$ coefficients. However, the main advantage of leptonic processes is that the MI are not greatly reduced from M_{GUT} to M_W. In this case, $(\delta^e_{LL})_{23} \simeq 2 \times 10^{-2}$ (except factors for order 1), contributing to $\tau \to \mu\gamma$ transitions, while the bound in Table 5-14 is only 3×10^{-2}. Therefore a $\tau \to \mu\gamma$ branching ratio close to the experimental bound is indeed possible.

Conclusions

We have introduced the mass insertion formalism and we have applied it to CP violation in B physics. We have seen that Super B Factoriescan explore the flavor structure of the sfermion mass matrices, both in the squark and in the slepton sectors. Supersymmetric Grand Unification predicts links between various leptonic and hadronic FCNC observables. We have quantitatively studied a SU(5) model and the implications for transitions between the second and third generations. We have shown that the present limit on $\mathcal{B}(\tau \to \mu\,\gamma)$ significantly constrains the observability of SUSY in CP violating B decays. In these models, lepton flavor-violating decays may be closer to experimental bounds than quark FCNCs, although these decays measure slightly different flavor parameters. We have also seen that sizable contributions are possible in "realistic" flavor models. Thus, precision measurements in B (and τ) physics are necessary to understand flavor physics.

5.3.5 Effective Supersymmetry in B Decays

$\succ$– P. Ko –$\prec$

Introduction

Generic SUSY models suffer from serious SUSY flavor and CP problems, because the squark mass matrices and quark mass matrices need not be simultaneously diagonalizable in the flavor space. Therefore the $\tilde{g} - \tilde{q}_{iA} - q_{jB}$ vertices ($i, j = 1, 2, 3$ are flavor indices, and $A, B = L, R$ denote chiralities) are described by some unitary matrix $W^d_{ij,AB}$ in the down (s)quark sector, which is analogous to the CKM matrix in the Standard Model. Since this coupling has a root in strong interaction and can have CP violating phases, it leads to too large flavor changing neutral current (FCNC) amplitudes through gluino-squark loop as well as too large ϵ_K and neutron electric dipole moment (EDM), which could easily dominate the Standard Model contributions and the data. Some examples are $K^0\overline{K}^0$ (ΔM_K and ϵ_K) and $B^0\overline{B}^0$ mixing (ΔM_B and CP asymmetry in $B_d \to J/\psi K_S^0$) and $B \to X_s\gamma$, *etc...* The lepton sector has the same problem through the neutralino-slepton loop, and the most serious constraint comes from $\mathcal{B}(\mu \to e\gamma)$ and electron EDM.

One way out of these SUSY flavor and CP problems is to assume that the first and second generation squarks are very heavy ($\gtrsim \mathcal{O}(10)$ TeV) and almost degenerate [206]. The third generation squarks and gauginos should be relatively light ($\lesssim 1$ TeV) in order that the quantum correction to Higgs mass parameter is still reasonably small. This scenario is called an effective SUSY model, or a decoupling scenario. In effective SUSY models, the $\tilde{b} - \tilde{g}$ loop can still induce a certain amount of flavor and CP violation in the quark sector through the mixing matrices $W^d_{ij,AB}$. In addition to the flavor mixing and CP violation from $W^d_{ij,AB}$'s, there could be flavor-conserving CP violation through the μ and A_t parameters within the effective SUSY models. Note that this class of models, ignoring the gluino-mediated FCNC, are used in the context of electroweak baryogenesis within SUSY (see, *e.g.*, [207]). Although the phases in μ and A_t are flavor-conserving, they can affect K and B physics through chargino/stop propagators and mixing angles.

Since there are no well-defined effective SUSY models as there are in gauge mediation or minimal supergravity scenarios, one has to assume that all the soft SUSY breaking terms have arbitrary CP-violating phases, as long as they satisfy the decoupling spectra and various experimental constraints. In order to make the analysis easy and transparent, we consider two extreme cases:

- Minimal flavor violation (MFV)
- Gluino-squark dominance in $b \to s(d)$ transition ($\tilde{g}$ dominance)

We will describe typical signatures of each scenario, keeping in mind that reality may involve a combination of these two extreme cases.

CP violation from μ and A_t phases

Let us first discuss minimal flavor violation models with effective SUSY spectra. In this model, flavor violation comes through the CKM matrix, whereas CP violation originates from the μ and A_t phases, as well as the CKM phase. The one loop electric dipole moment (EDM) constraint is evaded in the effective SUSY model due to the decoupling of the first/second generation sfermions, but there are potentially large two loop contribution to electron/neutron EDM's through Barr-Zee type diagrams in the large $\tan\beta$ region [208]. Imposing this two-loop EDM constraint and direct search limits on Higgs and SUSY particles, we find that [209, 210]

- There are no new phase shifts in $B^0\overline{B}^0$ and $B^0_s\overline{B}^0_s$ mixing: Time-dependent CP asymmetries in $B_d \to J/\psi K_S^0$ still measure the CKM angle $\beta = \phi_1$ [Fig. 5-36 (a)]
- ΔM_{B_d} can be enhanced up to $\sim 80\%$ compared to the Standard Model prediction [Fig. 5-36 (b)]
- Direct CP asymmetry in $B \to X_s\gamma$ ($A_{CP}^{b\to s\gamma}$) can be as large as $\pm 15\%$ [see Fig. 5-37]
- $R_{\mu\mu}$ can be as large as 1.8
- ϵ_K can differ from the Standard Model value by $\sim 40\%$

One can therefore anticipate substantial deviations in certain observables in the B system in SUSY models with minimal flavor violation and complex μ and A_t parameters. This class of models include electroweak baryogenesis (EWBGEN) within the MSSM and some of its extensions (such as NMSSM), where the chargino and stop sectors are the same as in the MSSM. In the EWBGEN scenario within the MSSM, the current lower limit on the Higgs mass requires a large radiative correction from the stop loop. Since $\tilde{t}_R$ has to be light to have a sufficiently strong 1st

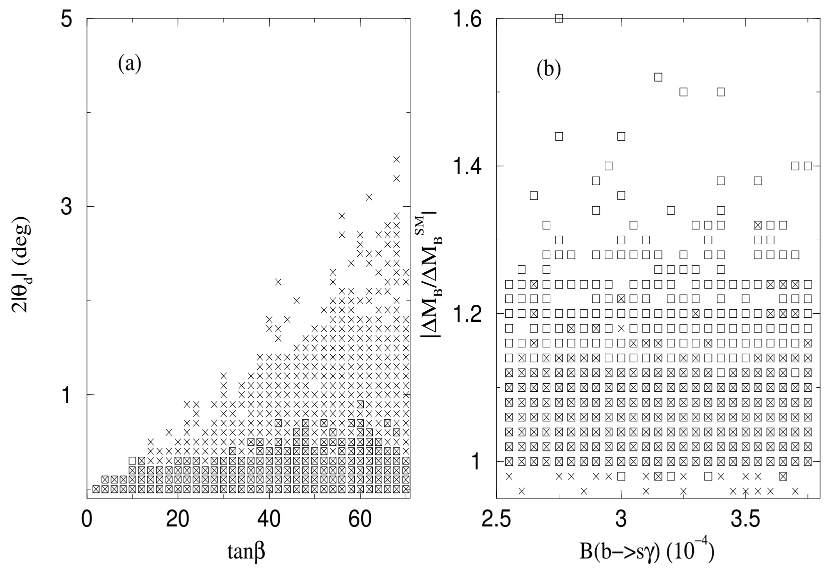

Figure 5-36. *Correlations between (a)* $\tan\beta$ *vs. the new phase shift in the* $BzBzb$ *mixing, and (b)* $\mathcal{B}(B \to X_s\gamma)$ *vs.* $|\Delta M_{B_d}/\Delta M_{B_d}^{\rm SM}|$*. The squares (the crosses) denote those which (do not) satisfy the two-loop EDM constraints.*

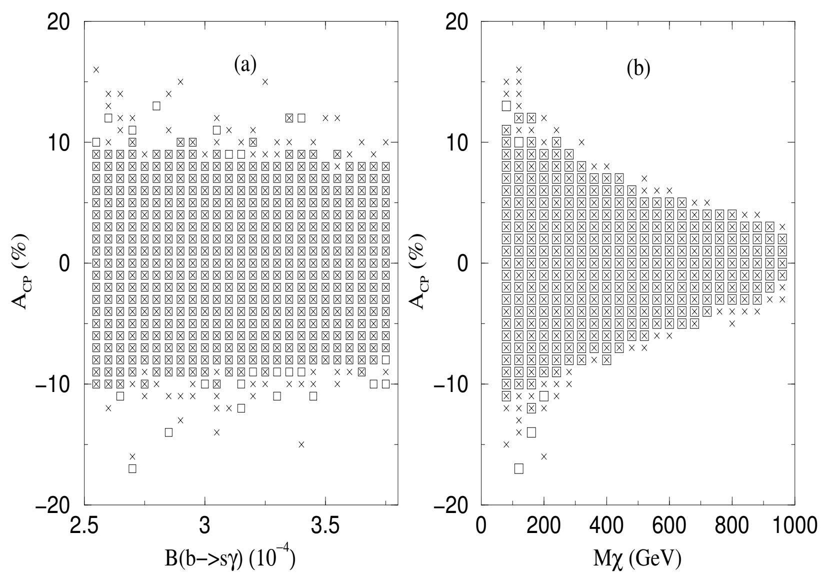

Figure 5-37. *Correlations of* $A_{CP}^{b\to s\gamma}$ *with (a)* $\mathcal{B}(B \to X_s\gamma)$ *and (b) the lighter chargino mass* $M_{\chi^\pm}$*. The squares (the crosses) denote those which (do not) satisfy the two-loop EDM constraints.*

order electroweak phase transition, one has to have heavy $\tilde{t}_L$ to induce a large Δm_h^2. After considering $B \to X_s\gamma$, one expects a very small deviations in $A_{CP}^{b\to s\gamma}$ and ΔM_{B_d} [211]. However, in some extensions of the MSSM, the tension between m_h and $m_{\tilde{t}_L}$ becomes significantly diluted in EWBGEN scenarios, because there could be tree level contributions to m_h^2. Therefore, the predictions made in Refs. [209, 210] will be still valid in EWBGEN scenarios beyond the MSSM.

Super B Factoriesshould be able to measure $A_{CP}^{b\to s\gamma}$ to higher accuracy, and will impose a strong constraint on a new CP-violating phase that could appear in $B \to X_s\gamma$. Also the forward-backward asymmetries in $B \to X_s\ell^+\ell^-$ with $\ell = e$ or μ are equally important probes of new CP-violating phases, and important observables to be measured at Super B Factories, for which LHCb or BTeV cannot compete.

CP violation from gluino-squark loops

In effective SUSY scenarios, it is possible that the gluino-mediated $b \to s$ transition is dominant over other SUSY contributions. Cohen *et al.* have described qualitative features of such scenarios in B physics [212]; a more quantitative analysis was presented by other groups. In Ref. [213], effects of possible new CP-violating phases on $B \to X_s\gamma$ and $B \to X_s\ell^+\ell^-$ were considered both in a model-independent manner, and in gluino mediation dominance scenario. In effective SUSY models, $A_{CP}^{b\to s\gamma}$ can be as large as $\pm 10\%$, if the third generation squarks are light enough $m_{\tilde{b}} \simeq (100-200)$ GeV (see Fig. 5-38), whereas $B \to X_s\ell^+\ell^-$ is almost the same as the Standard Model prediction [213].

How to distinguish the μ or A_t phase from the δ_{23}^d phase

Should we find deviations in $\sin 2\beta_{\phi K_S^0}$ or $A_{CP}^{b\to s\gamma}$, it will be very important to figure out the origin of new CP-violating phases. In an effective SUSY context, one has complex A_t, μ or $(\delta_{AB}^d)_{23}$ (with $A, B = L, R$). The effects of these new complex parameters on some oservables in the B system are shown in Table 5-16. The only process which is not directly affected by gluino-mediated FCNC is $B \to X_s\nu\overline{\nu}$. All the other observables are basically affected by both the phases of μ, A_t and $(\delta_{AB}^d)_{23}$ parameters. In fact, this feature is not specific to the effective SUSY scenarios, but is rather generic within SUSY models. Therefore the measurement of $B \to X_s\nu\overline{\nu}$ branching ratio will play a crucial role to tell if the observed CP-violating phenomena comes from the μ or A_t phase or $(\delta_{AB}^d)_{23}$. This can be done only at a e^+e^- Super B Factory, and not at hadron B factories.

Table 5-16. *Possible effects of the phase of μ or A_t for moderate $\tan\beta$ ($3 \le \tan\beta \le 6$) and the phase of δ_{i3}^d (with $i = 1, 2$) to various observables in the B systems, and possibilities to probe these at various experiments*

Observables	Arg (μ) or Arg (A_t)	Arg (δ_{i3}^d)	Super B Factory	LHCb
Δm_d	Y	Y	O	O
$\sin 2\beta$	N	Y	O	O
Δm_s	Y	Y	X	O
$\sin 2\beta_s$	N	Y	X	O
$A_{\rm CP}^{b\to s\gamma}$	Y	Y	O	X
$A_{\rm CP}^{b\to d\gamma}$	Y	Y	O	X
$B \to X_s\ell^+\ell^-$	Y	Y	O	X
$B \to X_s\nu\overline{\nu}$	Y	N	O	X
$B_d \to \phi K_S^0$	Y	Y	O	O

Conclusion

We showed that there could be large deviations in certain observables in the B system, which can be studied only in Super B Factories. The most prominent deviations are the branching ratio of $B \to X_s\nu\overline{\nu}$, $A_{CP}^{b\to s\gamma}$, the forward-backward asymmetry in $B \to X_s\ell^+\ell^-$, and the branching ratio of $B \to X_d\gamma$ and CP violation therein, without any

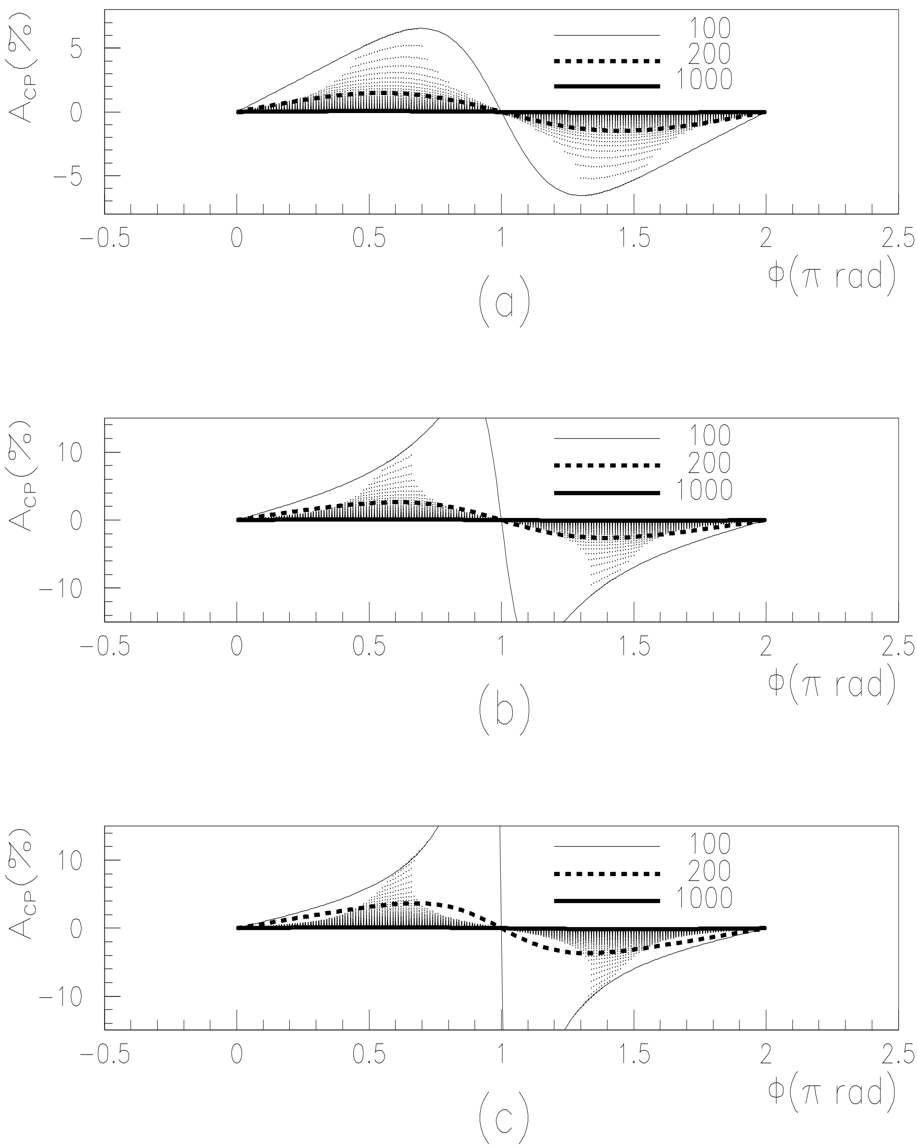

Figure 5-38. $A_{\rm CP}^{b\to s\gamma}$ *contours in the* $(\tilde{m},\phi)$ *plane for (a)* $x=0.3$*, (b)* $x=1$ *and (c)* $x=3$ *in the* (LL) *insertion case using the vertex mixing method with* $x=m_{\tilde{g}}^2/m_{\tilde{b}}^2$*.*

conflict with our current understanding based on the CKM paradigm. These observables could reveal new sources of CP and flavor violations that could originate from SUSY models, including effective SUSY models, and should be an important topic at Super B Factories.

I am grateful to S. Baek, Y. G. Kim and J. S. Lee for their collaboration on the work presented here.

5.3.6 Supersymmetric Flavor Violation: Higgs–Quark Interactions

$\succ$– D. A. Demir –$\prec$

The primary goal of the existing and planned hadron colliders and the B meson factories is to test the Standard Model and determine possible New Physics effects on its least understood sectors: breakdown of CP, flavor and gauge symmetries. In the standard picture, both CP and flavor violations are restricted to arise from the CKM matrix, and the gauge symmetry breaking is accomplished by introducing the Higgs field. However, the Higgs sector is badly behaved at quantum level; its stabilization against quadratic divergences requires supersymmetry (SUSY) or some other extension of the Standard Model. The soft breaking sector of the minimal SUSY model (MSSM) accommodates novel sources for CP and flavor violations [214, 215] with testable signatures at present (PEP-II, KEK-B) or future (Super B Factory or LHC) experiments. The Yukawa couplings, which are central to Higgs searches at the LHC, differ from all other couplings in one aspect: the radiative corrections from sparticle loops depend only on the ratio of the soft masses, and hence they do not decouple even if the SUSY-breaking scale lies far above the weak scale. In this sense, a non-standard hierarchy and texture of the Higgs-quark couplings, once confirmed experimentally, might provide direct access to sparticles, irrespective of how heavy they might be (though not too large to regenerate the gauge hierarchy problem). This section will summarize the results of recent work [216] that discusses the radiative corrections to Yukawa couplings from sparticle loops and their impact on flavor-changing neutral current (FCNC) observables and Higgs phenomenology.

The soft breaking sector mixes sfermions of different flavor via the off–diagonal entries of the sfermion mass–squared matrices. The LR and RL blocks are generated after the electroweak breaking with the maximal size $\mathcal{O}(m_t M_{SUSY})$, and their flavor-mixing potential is dictated by the Yukawa couplings $\mathbf{Y}_{u,d}$ and by the trilinear coupling matrices $\mathbf{Y}^A_{u,d}$ with $\left(\mathbf{Y}^A_{u,d}\right)_{ij} = \left(\mathbf{Y}_{u,d}\right)_{ij} \left(A_{u,d}\right)_{ij}$ where $A_{u,d}$ are not necessarily unitary, so that even their diagonal entries contribute to CP-violating observables. The flavor mixings in the LL and RR sectors, however, are insensitive to electroweak breaking; they are of pure SUSY origin. Clearly, CP violation in the LL and RR sectors is restricted to the flavor-violating entries, due to hermiticity. In discussing the FCNC transitions, it is useful to work with the mass insertions [215]

$$\left(\delta^{d,u}_{ij}\right)_{RR,LL} = \frac{\left(M^2_{D,U}\right)^{ij}_{RR,LL}}{\overline{M^2_{D,U}}}\,, \tag{5.155}$$

where $\left(M^2_D\right)_{RR,LL}$ have the generic form

$$\left(M^2_D\right)_{LL} = \begin{pmatrix} M^2_{\tilde{d}_L} & M^2_{\tilde{d}_L\tilde{s}_L} & M^2_{\tilde{d}_L\tilde{b}_L} \\ M^2_{\tilde{s}_L\tilde{d}_L} & M^2_{\tilde{s}_L} & M^2_{\tilde{s}_L\tilde{b}_L} \\ M^2_{\tilde{b}_L\tilde{d}_L} & M^2_{\tilde{b}_L\tilde{s}_L} & M^2_{\tilde{b}_L} \end{pmatrix} \quad , \quad \left(M^2_D\right)_{RR} = \begin{pmatrix} M^2_{\tilde{d}_R} & M^2_{\tilde{d}_R\tilde{s}_R} & M^2_{\tilde{d}_R\tilde{b}_R} \\ M^2_{\tilde{s}_R\tilde{d}_R} & M^2_{\tilde{s}_R} & M^2_{\tilde{s}_R\tilde{b}_R} \\ M^2_{\tilde{b}_R\tilde{d}_R} & M^2_{\tilde{b}_R\tilde{s}_R} & M^2_{\tilde{b}_R} \end{pmatrix} \tag{5.156}$$

in the bases $\{\tilde{d}_L, \tilde{s}_L, \tilde{b}_L\}$ and $\{\tilde{d}_R, \tilde{s}_R, \tilde{b}_R\}$, respectively. The same structure repeats for the up sector. The mass insertions are defined in terms of $\overline{M^2_{D,U}}$ which stand for the mean of diagonal entries. The textures of the LL and RR blocks are dictated by the SUSY breaking pattern. In minimal SUGRA and its nonuniversal variants with CP

violation, for instance, the size and structure of flavor and CP violation are dictated by the CKM matrix [214]. On the other hand, in SUSY GUTs with Yukawa unification *e.g.*, SO(10), implementation of the see–saw mechanism for neutrino masses implies sizable flavor violation in the RR block, given the large mixing observed in atmospheric neutrino data [217].

The effective theory below the SUSY breaking scale M_{SUSY} consists of a modified Higgs sector; in particular, the tree level Yukawa couplings receive sizable corrections from sparticle loops [216]. For instance, the d quark Yukawa coupling relates to the physical Yukawas via

$$h_d = \frac{g_2 \overline{m}_d}{\sqrt{2} M_W \cos\beta} \, \frac{1 - a^2 \left(\delta^d_{23}\right)_{LR} \left(\delta^d_{32}\right)_{LR} - a A_{12} \, \frac{\overline{m}_s}{\overline{m}_d} - a A_{13} \, \frac{\overline{m}_b}{\overline{m}_d}}{1 - a^2 A_2 - a^3 A_3} \,, \tag{5.157}$$

where $a = \epsilon \tan\beta / (1 + \epsilon \tan\beta)$, $A_{12} = \left[\left(\delta^d_{12}\right)_{LR} - a \left(\delta^d_{13}\right)_{LR} \left(\delta^d_{32}\right)_{LR} \right]$, $A_{13} = A_{12}(2 \leftrightarrow 3)$, $A_2 = \left| \left(\delta^d_{12}\right)_{LR} \right|^2 + \left| \left(\delta^d_{13}\right)_{LR} \right|^2 + \left| \left(\delta^d_{23}\right)_{LR} \right|^2$ and $A_3 = \left(\delta^d_{12}\right)_{LR} \left(\delta^d_{23}\right)_{LR} \left(\delta^d_{31}\right)_{LR} +$ h.c. Here $\epsilon = (\alpha_s / 3\pi) e^{-i(\theta_\mu + \theta_g)}$, and

$$\left(\delta^d_{ij}\right)_{LR} = \frac{1}{6} \, \left(\delta^d_{ij}\right)_{RR} \left(\delta^d_{ji}\right)_{LL} \,, \tag{5.158}$$

with the SUSY CP-odd phases defined as $\theta_g = \mathrm{Arg}[M_g]$, $\theta_\mu = \mathrm{Arg}[\mu]$, $\theta^d_{ij} = \mathrm{Arg}[(A_d)_{ij}$, *etc.*As (5.157) suggests, in contrast to the minimal flavor violation (MFV) scheme, the Yukawa couplings acquire large corrections from those of the heavier ones. Indeed, the radiative corrections to $h_d/\overline{h}_d$, $h_s/\overline{h}_s$, $h_u/\overline{h}_u$ and $h_c/\overline{h}_c$ involve, respectively, the large factors $\overline{m}_b/\overline{m}_d \sim (\tan\beta)^2_{max}$, $\overline{m}_b/\overline{m}_s \sim (\tan\beta)_{max}$, $\overline{m}_t/\overline{m}_u \sim (\tan\beta)^3_{max}$, and $\overline{m}_t/\overline{m}_c \sim (\tan\beta)^2_{max}$ with $(\tan\beta)_{max} \lesssim \overline{m}_t/\overline{m}_b$. Unlike the light quarks, the top and bottom Yukawas remain close to their MFV values, to a good approximation. Therefore, the SUSY flavor-violation sources mainly influence the light quark sector, thereby modifying several processes in which they participate. These corrections are important even at low $\tan\beta$. As an example, consider $\left(\delta^d_{13}\right)_{LR} \sim 10^{-2}$ for which $h_d/h_d^{MFV} \simeq 0.02(2.11), -2.3(6.6), -4.6(17.7)$ for $\tan\beta = 5, 20, 40$ at $\theta_\mu + \theta_g \to 0(\pi)$. Note that the Yukawas are enhanced especially for $\theta_\mu + \theta_g \to \pi$, which is the point preferred by Yukawa–unified models such as SO(10). In general, as $\tan\beta \to (\tan\beta)_{max}$ the Yukawa couplings of down and strange quarks become amproximately degenerate with the bottom Yukawa for $\left(\delta^d_{13,23}\right)_{LR} \sim 0.1$ and $\theta_\mu + \theta_g \to \pi$. There is no $\tan\beta$ enhancement for up quark sector but still the large ratio $\overline{m}_t/\overline{m}_u$ sizably folds h_u compared to its Standard Model value: $h_u \simeq 0.6 \, e^{i(\theta^u_{11} - \theta_g)} \, \overline{h}_c$ for $(\delta^u_{13})_{LR} \sim 0.1$.

The SUSY flavor violation influences the Higgs-quark interactions by (i) modifying $H^a \overline{q} q$ couplings via sizable changes in Yukawa couplings as in (5.157), and by (ii) inducing large flavor changing couplings $H^a \overline{q} q'$:

$$\begin{aligned} & \frac{\overline{h}_{d^i}^{\;SM}}{\sqrt{2}} \left[\frac{h^i_d}{\overline{h}^i_d} \, \tan\beta \, C^d_a + \left(\frac{h^i_d}{\overline{h}^i_d} - 1 \right) \left(e^{i(\theta^d_{ii} + \theta_\mu)} \, C^d_a - C^{u\star}_a \right) \right] \overline{d}^i_R \, d^i_L \, H_a \\ + \; & \frac{\overline{h}^{SM}_{d^i}}{3\sqrt{2}} \epsilon \tan\beta \left[\frac{h^i_d}{\overline{h}_{d^i}} \, \left(\delta^d_{ij}\right)_{LL} + \frac{h^j_d}{\overline{h}^i_d} \, \left(\delta^d_{ij}\right)_{RR} \right] \left(\tan\beta \, C^d_a - C^{u\star}_a \right) \overline{d}^i_R \, d^j_L \, H_a \end{aligned} \tag{5.159}$$

where $C^d_a \equiv \{-\sin\alpha, \cos\alpha, i \sin\beta, -i \cos\beta\}$ and $C^u_a \equiv \{\cos\alpha, \sin\alpha, i\cos\beta, i \sin\beta\}$ in the basis $H_a \equiv \{h, H, A, G\}$ if the CP violation effects in the Higgs sector, which can be quite sizable [218] and add additional CP-odd phases [219] to Higgs-quark interactions, are neglected . Similar structures also hold for the up sector. The interactions contained in (5.159) have important implications for both FCNC transitions and Higgs decay modes. The FCNC processes are contributed by both the sparticle loops (*e.g.*, the gluino-squark box diagram for $K^0 \overline{K}^0$ mixing) and Higgs exchange amplitudes. The constraints on various mass insertions can be satisfied by a partial cancellation between these two contributions if M_{SUSY} is close to the weak scale. On the other hand, if M_{SUSY} is high, then the only surviving SUSY contribution is the Higgs exchange. In either of these extremes, or in-between, the main issue is to determine what size and phase the FCNC observables allow for the mass insertions. This certainly requires a global analysis of the existing FCNC data by incorporating the Higgs exchange effects to other SUSY contributions [220]. For example, in parameter regions where the latter are suppressed ($M_{SUSY} \gg m_t$), one can determine the allowed sizes of mass insertions by using (5.157) in (5.159). Doing so, one finds that the flavor–changing Higgs vertices bsH^a and bdH^a become vanishingly small for $\tan\beta \simeq 60$ when all MIs are $\mathcal{O}(1)$, for $\tan\beta \simeq 65$ when $\left(\delta^d_{12}\right)_{LL,RR} \simeq 0$, and, finally, for $\tan\beta \simeq 68$,when $\left(\delta^d_{12}\right)_{LL} \simeq -\left(\delta^d_{12}\right)_{RR}$, provided that $\phi_\mu + \phi_g \to \pi$ in all three cases. Therefore, in this parameter domain, though the flavor-changing Higgs decay channels are sealed up, the decays into similar quarks are highly enhanced. For instance, $\Gamma(h \to \overline{d}d)/\Gamma(h \to \overline{b}b) \simeq (\text{Re}\,[h_d/h_b])^2$ which is $\mathcal{O}(1)$ when $h_d \sim h_b$, as is the case with SUSY flavor violation. Such enhancements in light quark Yukawas induce significant reductions in $\overline{b}b$ branching fractions — which is a very important signal for hadron colliders to determine the non-standard nature of the Higgs boson ($h \to \overline{b}b$ has $\sim 90\%$ branching fraction in the Standard Model). If FCNC constraints are saturated without a strong suppression of the flavor-changing Higgs couplings (which requires M_{SUSY} to be close to the weak scale) then Higgs decays into dissimilar quarks get significantly enhanced. For instance, $h \to \overline{b}s + \overline{s}b$ can be comparable to $h \to \overline{b}b$. (See [221] for a diagrammatic analysis of $\to \overline{b}s + \overline{s}b$ decay.) In conclusion, as fully detailed in [216], SUSY flavor and CP violation sources significantly modify Higgs–quark interactions, thereby inducing potentially large effects that can be discovered at Super B Factories, as well as at the hadron colliders.

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5.4 Extra Dimensions

5.4.1 Large Extra Dimensions and Graviton Exchange in $b \to s\ell^+\ell^-$

⋗ T. G. Rizzo ⋖

Introduction

The existence of extra space-like dimensions has been proposed as a possible solution to the gauge hierarchy problem. Although there are many models in the literature attempting to address this issue a common feature is the existence of a higher dimensional 'bulk' space in which gravity is free to propagate. In the Kaluza-Klein (KK) picture the reduction to four dimensions leads to the existence of a massive tower of gravitons that can be exchanged between Standard Model fields. The two most popular scenarios are those of Arkani-Hamed, Dimopolous and Dvali (ADD)[222] and of Randall and Sundrum (RS)[223]. The properies of the KK gravitons are significantly different in these two models. However, in either case the exchange of KK gravitons has been shown to lead to unique signatures that may be discovered at the LHC[224].

While high-p_T measurements at hadron colliders may tell us some of the gross features of the extra-dimensional model, other sets of measurements will be necessary in order to disentangle its complete structure. For example, the LHC may observe the graviton resonances of the RS model in the Drell-Yan and/or dijet channels, but it will be very difficult, if not impossible, to examine the possible *flavor* structure of graviton couplings in such an enviornment[225]. While such determinations will certainly be possible at a future Linear Collider, provided it has suffient center-of-mass energy to sit on a graviton resonance, it may be a while between the LHC discovery and the data from the Linear Collider becoming available. It is possible, however, that at least some aspects of the flavor structure of the graviton KK couplings may be determined using precision data at lower energies, through rare decays such as $b \to s\ell^+\ell^-$. This is the subject of the discussion below.

Flavor dependence, as well as flavor violation in KK graviton couplings can be generated in models which attempt to explain the fermion mass hierarchy as well as the structure of the CKM matrix[226]. In such scenarios, fermions are localized in the extra dimensions either via scalar 'kink'-like solutions, or via their 5-d Dirac masses. A description of the details of such models is, however, beyond the scope of this discussion. In fact, wishing to be as model-independent as possible, we note that in all scenarios at low energies the exhange of gravitons between Standard Model fields can be described by the single dimension-8 operator

$$O_{grav} = \frac{1}{M^4}\, X\, T_{\mu\nu} T^{\mu\nu}\,, \tag{5.160}$$

where M is a mass scale of order $\sim$ a few TeV, the $T_{\mu\nu}$ are the stress-energy tensors of the Standard Model fields, which can have complex flavor structures, and X is a general coupling matrix. Operators such as these may be generated in either ADD-like or RS-like scenarios but we will not be interested here in the specific model details. Instead we focus on unique signatures for graviton exchange associated with the above operator.

Analysis

How can $b \to s\ell^+\ell^-$ probe such operators? To be specific, let us consider the case of ADD-like models; as we will see, our results are easily generalized to RS-type scenarios. In ADD, we identify $M \to M_H$, the cutoff scale in the theory, and $X \to \lambda X$, with λ being an overall sign. Identifying the first(second) $T_{\mu\nu}$ with the $b\bar{s}(\ell^+\ell^-)$-effective graviton vertex, the new operator will lead to, *e.g.*, a modification of the $b \to s\ell^+\ell^-$ differential decay distribution. Following the notation in [227], we find that this is now given by

$$\begin{aligned}\frac{d^2\Gamma}{dsdz} \sim\ & [(C_9 + 2C_7/s)^2 + C_{10}^2][(1+s) - (1-s)z^2] - 2C_{10}(C_9 + 2C_7/s)sz\,,\\ &+ \frac{4}{s^2}C_7^2(1-s)^2(1-z^2) - \frac{4}{s}C_7(C_9 + 2C_7/s)(1-s)(1-z^2)\,,\end{aligned}$$

$$+ DC_9(1-s)z[2s+(1-s)z^2] + DC_{10}s(1-s)(1-z^2)\,, \tag{5.161}$$

where the C_i are the usual effective Standard Model Wilson coefficients, $s = q^2/m_b^2$ is the scaled momentum transfer, $z = \cos\theta$ is the dilepton pair decay angle, and

$$D = \frac{2m_b^2}{G_F\alpha}\sqrt{2}\pi\frac{1}{V_{tb}V_{ts}}\frac{\lambda X}{M_H^4} \simeq 0.062\frac{\lambda X}{M_H^4} \tag{5.162}$$

describes the strength of the graviton contribution with M_H in TeV units. The terms proportional to D in this expression result from the interference of the Standard Model and graviton KK tower exchange amplitudes; note that there is no term proportional to DC_7, as dipole and graviton exchanges do not interfere. Here we neglect the square of the pure graviton contribution in the rate, since it is expected to be small.

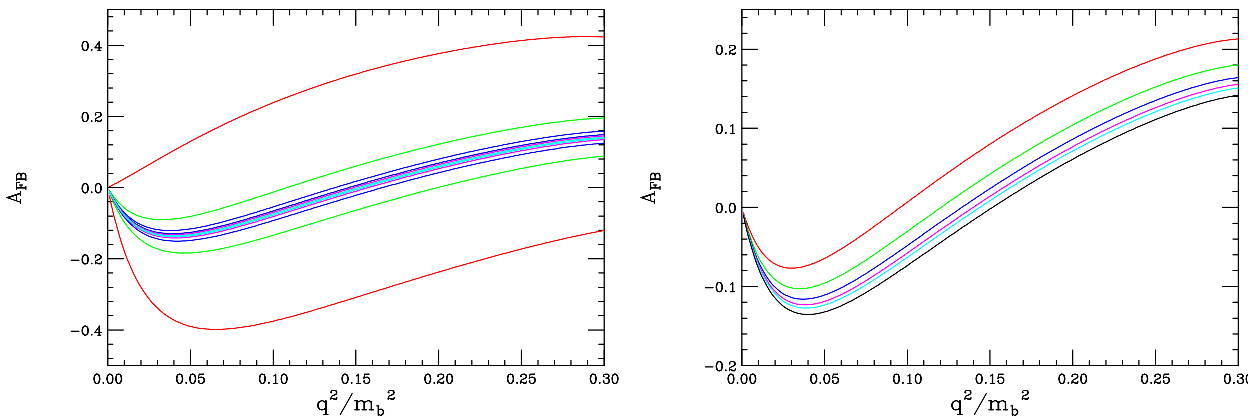

Figure 5-39. *A_{FB} as a function of s in the ADD(left) and RS(right) scenarios. In the ADD case, from outside to inside the curves are for $M_H = 1, 1.5, 2.....$ TeV with results for both signs of λ shown. In the RS case, from left to right the curves correspond to masses of the lightest KK graviton being 600, 700,...GeV with $k/M_{pl} = 0.1$ being assumed. In either case we take $X = 1$ for purposes of demonstration; for other X values the curves will scale as $M \to M/X^{1/4}$. The current collider bounds correspond to $M_H > 1$ TeV and $m_1 > 600$ GeV, respectively.*

How can graviton exchange be observed uniquely in $b \to s\ell^+\ell^-$? As is well known, many sources of New Physics can lead to modifications in the $b \to s\ell^+\ell^-$ differential distribution[228]. In particular, one quantity of interest is the forward-backward asymmetry A_{FB} and the location of its corresponding zero as a function of s[229]. That graviton KK tower exchange modifies the location of the zero is clear from the expression above. Fig. 5-39 shows the typical shifts in A_{FB} and its zero in both ADD- and RS-like scenarios. Clearly any observable shifts due to graviton exchange are not by any means unique though they are signatures for New Physics.

Graviton exchange *does*, however, lead to a new effect which will be absent in all other cases of New Physics. The source of this new distinct signature is the z^3 term in the differential distribution above, which can be traced back to the spin-2 nature of graviton exchange. The existence of this type of term can be observed experimentally by using the moment method[230] previously employed to probe for KK graviton tower exchange in fermion pair production at the Linear Collider. To this end, we define the quantity

$$< P_3(s) >= \frac{\int \frac{d^2\Gamma}{dsdz}P_3(z)\,dz}{\frac{d\Gamma}{ds}} \tag{5.163}$$

where $P_3 = z(5z^2-3)/2$ is the third Legendre polynomial. Due to the orthogonality of the P_n, the presence of the z^3 term induces a non-zero value for this moment; the terms that go as $\sim z^{0,1,2}$ in the distribution yield zero for this observable. *Any* experimental observation of a non-zero value for this moment would signal the existence of flavor-changing gravitational interactions. Fig. 5-40 shows the typical $s-$dependence of this moment in both the ADD-like and RS-like scenarios. It now becomes an experimental issue as to whether or not such a non-zero moment

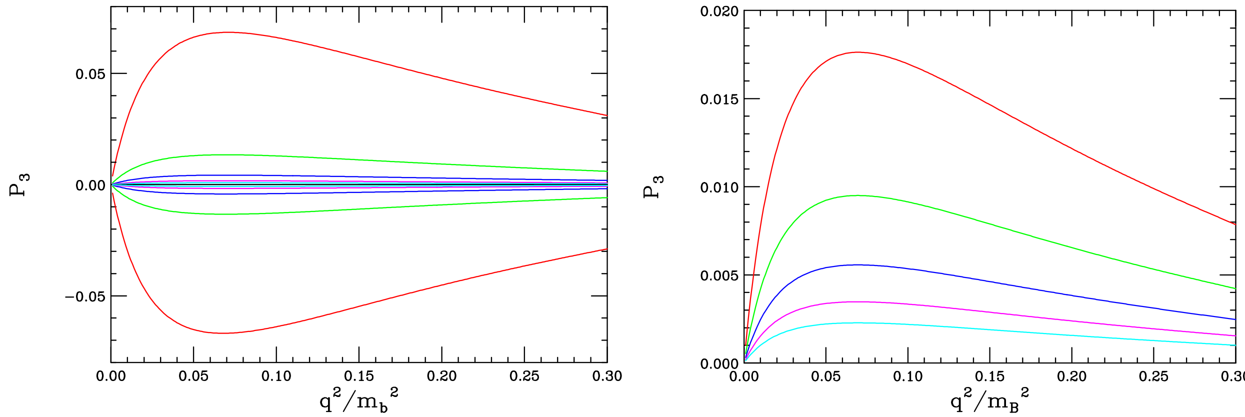

Figure 5-40. *Same as the previous figure but now showing the quantity $< P_3 >$ as a function of s.*

is observable. Clearly a very large statistical sample will be required on the order of $\sim 50 - 100\ \mathrm{ab}^{-1}$ or so. To reach this level, a Super B Factory is required. An experimental simulation along these lines would be useful.

In conclusion, we have shown that flavor-changing KK graviton exchange can be probed via the $b \to s\ell^+\ell^-$decay. A unique signature for these contributions can be obtained through the use of the moment technique. A nonzero value of the third Legendre moment will prove the existence of spin-2 exchange in this process. A Super B Factory is needed to reach the required level of statistics.

The author would like to thank J. L. Hewett for discussions related to this work.

5.4.2 TeV^{-1}-sized Extra Dimensions with Split Fermions

>– B. Lillie –<

Extra dimensions, in addition to the virtues already discussed, also present the possibility of understanding geometrically several dimensionless numbers that are observed to be very small. These include the small rate of proton decay, and the large ratios of fermion masses. This was first noted by Arkani-Hamed and Schmaltz [231]. They noticed that if the zero modes of the fermion fields were localized to Gaussians in the extra dimensions, then effective 4-d operators that contain fermions will be proportional to the overlaps of these Gaussians. If the localized fermions are separated from each other, these overlap integrals can be exponentially small. For example, separating quark and lepton fields by a distance a in one extra dimension (Fig. 5-41) results in a suppression of the proton decay operator $qqq\ell$ by

$$\int_0^R dy e^{-3\frac{y^2}{\sigma^2}} e^{-\frac{(y-a)^2}{\sigma^2}} = e^{-\frac{3}{4}\frac{a^2}{\sigma^2}} \tag{5.164}$$

where σ is the width of the fermions.

If the Higgs field lives in the bulk, then the fermion masses are generated by the flat zero mode of the Higgs, and are proportional to the overlap of the left- and right-handed fields. If the chiral components of different fermions are separated by different distances in the extra dimension, then exponentially different masses can be generated. The Yukawa coupling between the i-th left handed and j-th right-handed fermions is proportional to

$$\int_0^R dy e^{-\frac{(y-y_i)^2}{\sigma^2}} e^{-\frac{(y-y_j)^2}{\sigma^2}} = e^{-\frac{1}{2}\frac{(y_i-y_j)^2}{\sigma^2}} \ . \tag{5.165}$$

Thus, if this scenario were true, we could understand the large ratios of fermion masses as being due to order one differences in the parameters of the fundamental theory. It has been shown by explicit construction that the observed

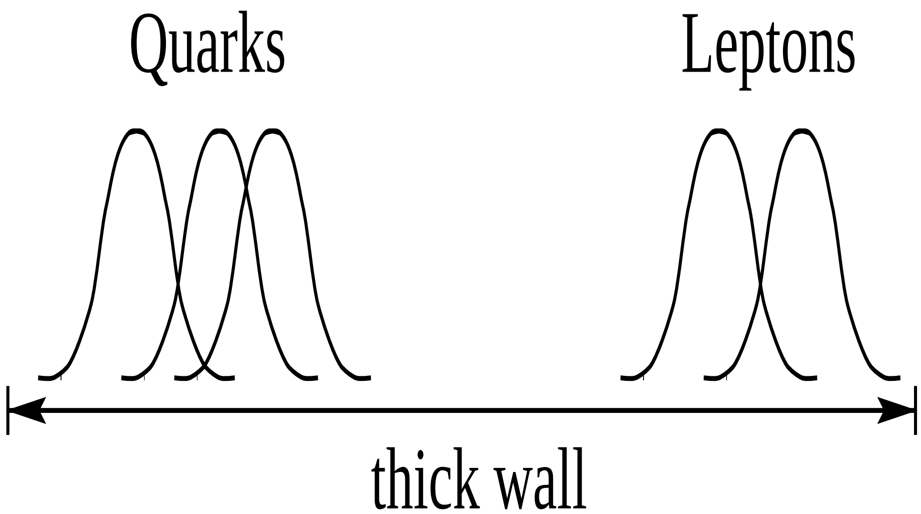

Figure 5-41. *Illustration of the concept of Split Fermions. Different species of fermions are localized to Gaussians at different locations in an extra dimension. This can be interpreted as a dimension compactified at the* TeV *scale, or a "brane" of* TeV *thickness embedded in a larger extra dimension. Figure taken from [231].*

values of the fermion masses (the Yukawa hierarchies), as well as the CKM matrix elements can be obtained in this way [232]. In addition, several variant models have been proposed. The left and right-handed fermions could be exponentially localized to two different branes, and the Higgs field to the left-handed brane. The Yukawa hierarchies are then obtained from the exponentially small values of the right-handed fermion wavefunctions at the left-handed brane [233]. This scenario generalizes very nicely to the case of a warped extra dimension, where the Higgs is localized to the TeV brane and all the fermions, except the top, are localized near the Plank brane [234]. Finally, rather than fixed-width Gaussians separated by some distance, one could consider different width Gaussians localized to the same point. Instead of exponentially small Yukawa matrix elements, this scenario generates Yukawa matrix elements that are all approximately the same, realizing the democratic scenario of fermion masses [235].

Split fermion scenarios naturally suppress many dangerous operators, but they do not suppress flavor changing effects [236, 237]. It is for this reason that they are of interest to a Super B Factory. To see this, note that, while fermions can be localized to different points in an extra dimension, the gauge fields must interact universally with the matter fields, and hence must posses a flat (or nearly flat) zero mode. This implies that they are delocalized at least on the scale of the separation of the fermions. They will then have excited KK modes with non-trivial wavefunctions. These will interact non-universally with the matter fields, and hence will produce flavor-changing effects. This is due to the fact that the non-universal couplings pick out a direction in flavor space, so that when the Yukawa matrices are diagonalized to find the mass eigenstates, non-trivial effects will be seen in the KK couplings. In the case of a single, flat, extra dimension, the coupling of the n-th excited gauge boson to a fermion localized to the point ℓ is proportional to

$$\int_0^R dy\ \cos(n\pi y) e^{-(y-\ell)^2 R^2/\sigma^2} \approx \cos(n\pi\ell) e^{-n^2\sigma^2/R^2}. \quad (5.166)$$

All excited gauge bosons, including the excited gluons, will have non-universal couplings and can generate flavor-changing neutral currents at tree-level. Hence, the model contains tree-level FCNC effects, suppressed only by the mass of the KK excitations. These can produce large effects, and thus already produce strong constraints from existing data.

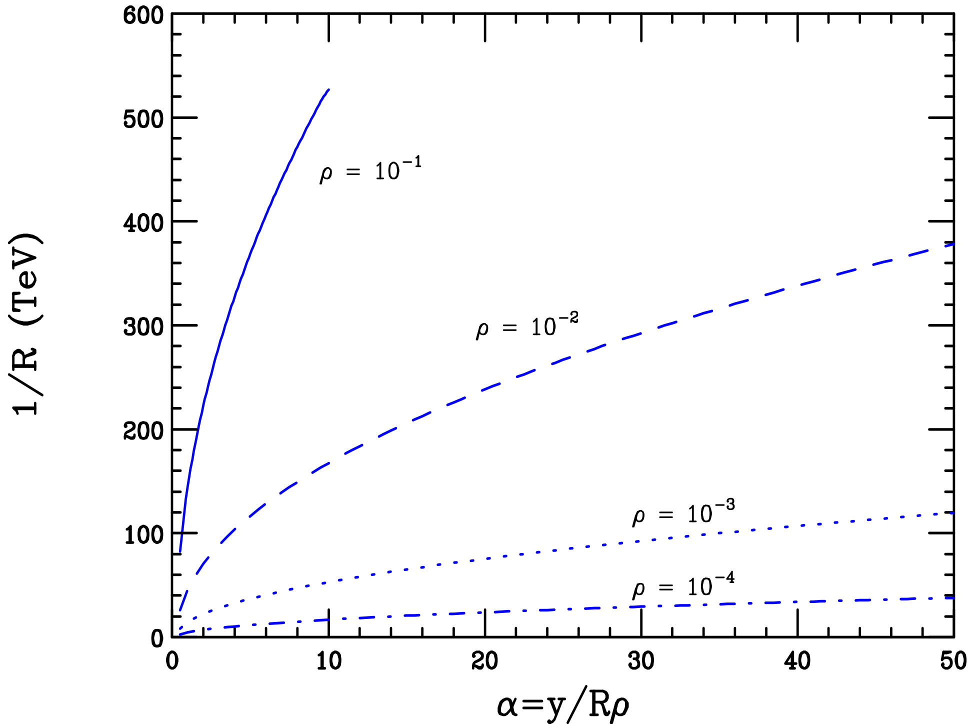

Figure 5-42. *Constraints on the compactification scale $1/R$ arising from the FCNC contribution to Δm_K, as a function of the separation of the fermions in units of the fermion width. The regions below the curves are excluded. Different curves are different values of $\rho = \sigma/R$, the ratio of the compactification scale to the fermion localization scale. Note that the size of the extra dimension in units of the fermion width is $1/\rho$, so for $\rho = 1/10$, $\alpha = 10$ corresponds to the fermions being localized at opposite ends of the dimension.*

Since all KK states contribute, any FCNC effect must be summed over all states. In one flat dimension, this sum is approximated by the distance between the two relevant flavors, in units of the size of the dimension. Note that the size of Yukawa matrix elements depends only on the separations in units of the fermion width, σ, and is independent of R. Hence, if we are interested in the flavor-changing effects in a scenario where the fermion masses are explained by localization, the relevant parameter is $\rho = \sigma/R$.

For example, the contribution to the mass splitting of an oscillating neutral meson system, P, from the separation of a pair of split fermions, takes the form

$$\Delta m_P = \frac{2}{9} g_s^2 f_P^2 m_P R^2 V^4_{qq',qq'} F(\rho\alpha). \tag{5.167}$$

Here $F(\rho\alpha)$ is a function that depends on the extra dimensional splitting, α of the two quarks, q and q', that form the meson, and V^4 represents four mixing angles arising from the matrices that diagonalize the Yukawa matrices (each one is roughly the square root of a CKM matrix element). The overall scale is set by the factor of R^2, and hence this formula can be interpreted as a constraint on the compactification scale from the measured value of Δm_P. Fig. 5-42 shows this constraint from Δm_K for several values of ρ. Note that extremely high scales can be probed.

Split fermion models are expected to live within a larger model that solves the hierarchy problem, and hence the cutoff of the theory is expected to be not much larger than 10 TeV. The compactification scale must be even smaller. We see that one can achieve a small value of $1/R$ at the expense of going to a very small value of ρ. However, this implies that the fermion localization scale $1/\sigma$ is very large, of order 10^4 TeV. Hence it looks like the models are essentially ruled out. This conclusion can be escaped in two ways. The bounds shown are for the kaon mass splitting,

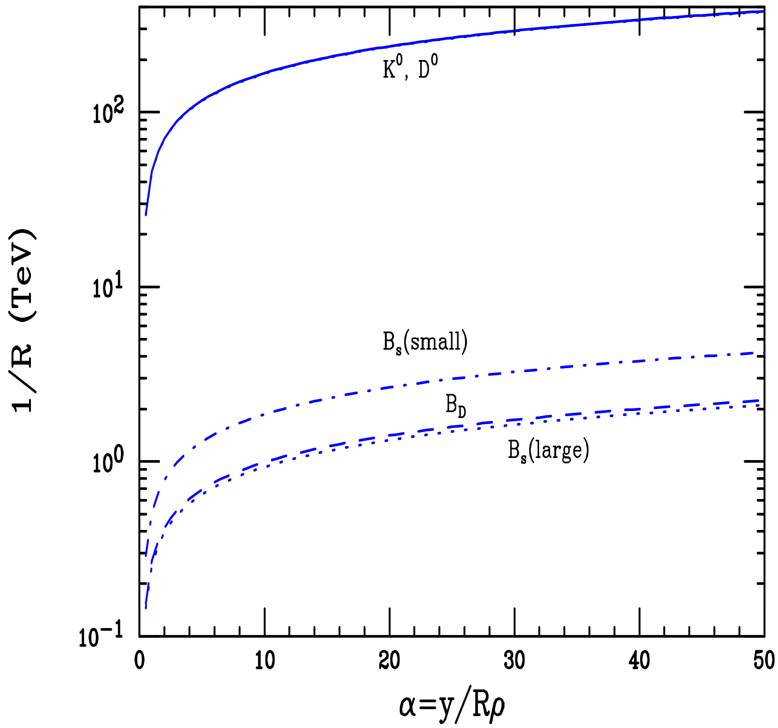

Figure 5-43. *Constraints on the compactification scale arising from the mass splitting of different neutral meson systems for $\rho = 1/100$. For B_s, since no upper bound is known, two different values were taken, the "small" value is about the size expected in the Standard Model, the "large" value about four times bigger. A combination of all of these measurement, and others, is needed to fully constrain the models.*

and hence are most sensitive to the separation of the d and s quarks. It could be that the down-type quarks are all aligned, and the CKM mixing is induced by splittings of the up-type quarks. To constrain this, one would need to look at $D^0 - \overline{D}^0$ mixing. Alternately, the mixing could be in both the up and the down sector, but such that the Yukawa matrices are diagonalized by mixing that is predominantly between the third and first or second generations, in such a way as to produce the CKM matrix when the product of up and down diagonalization matrices is taken. Constraining this requires measurement of both the B_d and B_s mixing parameters. Fig. 5-43 shows the constraints from all neutral meson mixings. A combination of these is needed to fully constrain the model.

It is also possible that there may be interesting signatures in rare decays. In particular, lepton family number-violating decays can produce limits on the splittings of leptons, which are not available from the meson oscillation data. It is also possible in a small region of parameter space to have contributions to $B^0 \to J/\psi K^0_S$ or $B^0 \to \phi K^0_S$ that are near the same order as the Standard Model, leading to interesting effects in CP-violating observables.

Another, more attractive, possibility is to go from a flat to a warped geometry. In [234] it was shown that the same flavor constraints are much more mild in an RS model where the fermions live in the bulk, but are localized near the Plank brane. In that case, the fact that the gauge KK wavefunctions are nearly flat near the Plank brane helps to naturally suppress the non-universality of the couplings to fermions. As a result all scales in the model can be below $10\ \mathrm{TeV}$, but there are still effects predicted in rare decays that might be visible to future experiments. This case is the most promising for future study at a Super B Factory.

5.4.3 Universal Extra Dimensions

≻— A. J. Buras, A. Poschenrieder, M. Spranger, and A. Weiler —≺

Introduction

Models with more than three spatial dimensions have been used to unify the forces of nature ever since the seminal papers of Kaluza and Klein [238]. More recently, extra dimensional models have been employed as an alternative explanation of the origin of the TeV scale [239].

A simple model of the so called universal type is the Appelquist, Cheng and Dobrescu (ACD) model [240] with one universal extra dimension. It is an extension of the Standard Model to a 5 dimensional orbifold $\mathcal{M}^4 \times S^1/Z_2$, where all the Standard Model fields live in all available 5 dimensions. In what follows we will briefly describe this model and subsequently report on the results of two papers [241, 83] relevant for these Proceedings, in which we investigated the impact of the KK modes on FCNC processes in this model. Further details can be found in Ref. [242].

The ACD Model

The full Lagrangian of this model includes both the boundary and the bulk Lagrangian. The coefficients of the boundary terms, although volume-suppressed, are free parameters and will get renormalized by bulk interactions. Flavor non-universal boundary terms would lead to large FCNCs. In analogy to a common practice in the MSSM, in which the soft supersymmetry breaking couplings are chosen to be flavor-universal, we assume negligible boundary terms at the cut-off scale. Now the bulk Lagrangian is determined by the Standard Model parameters, after an appropriate rescaling. With this choice, contributions from boundary terms are of higher order, and we only have to consider the bulk Lagrangian for the calculation of the impact of the ACD model.

Since all our calculations are cut-off-independent (see below) the only additional free parameter relative to the Standard Model is the compactification scale $1/R$. Thus, all the tree level masses of the KK particles and their interactions among themselves and with the Standard Model particles are described in terms of $1/R$ and the parameters of the Standard Model. This economy in new parameters should be contrasted with supersymmetric theories and models with an extended Higgs sector. All Feynman rules necessary for the evaluation of FCNC processes can be found in [241, 83].

A very important property of the ACD model is the conservation of KK parity that implies the absence of tree level KK contributions to low energy processes taking place at scales $\mu \ll 1/R$. In this context the flavor-changing neutral current(FCNC) processes like particle-antiparticle mixing, rare K and B decays and radiative decays are of particular interest. Since these processes first appear at one-loop in the Standard Model and are strongly suppressed, the one-loop contributions from the KK modes to them could in principle be important.

The effects of the KK modes on various processes of interest have been investigated in a number of papers. In [240, 243] their impact on the precision electroweak observables assuming a light Higgs ($m_H \leq 250$ GeV) and a heavy Higgs led to the lower bound $1/R \geq 300\,\mathrm{GeV}$ and $1/R \geq 250\,\mathrm{GeV}$, respectively. Subsequent analyses of the anomalous magnetic moment [244] and the $Z \to b\bar{b}$ vertex [245] have shown the consistency of the ACD model with the data for $1/R \geq 300\,\mathrm{GeV}$. The latter calculation has been confirmed in [241]. The scale of $1/R$ as low as $300\,\mathrm{GeV}$ would also lead to an exciting phenomenology in the next generation of colliders and could be of interest in connection with dark matter searches. The relevant references are given in [83].

The question then arises whether such low compactification scales are still consistent with the data on FCNC processes. This question has been addressed in detail in [241, 83]. Before presenting the relvant results of these papers let us recall the particle content of the ACD model that has been described in detail in [241].

In the effective four dimensional theory, in addition to the ordinary particles of the Standard Model, denoted as zero $(n = 0)$ modes, there are infinite towers of the KK modes $(n \geq 1)$. There is one such tower for each Standard Model boson and two for each Standard Model fermion, while there also exist physical neutral ($a^0_{(n)}$) and charged ($a^{\pm}_{(n)}$)

scalars with $(n \geq 1)$ that do not have any zero mode partners. The masses of the KK particles are universally given by

$$(m_{(n)}^2)_{\rm KK} = m_0^2 + \frac{n^2}{R^2} \,. \tag{5.168}$$

Here m_0 is the mass of the zero mode, as M_W, M_Z, m_t respectively. For $a_{(n)}^0$ and $a_{(n)}^\pm$ this is M_Z and M_W, respectively. In phenomenological applications it is more useful to work with the variables x_t and x_n defined through

$$x_t = \frac{m_t^2}{m_W^2}, \qquad x_n = \frac{m_n^2}{m_W^2}, \qquad m_n = \frac{n}{R} \tag{5.169}$$

than with the masses in (5.168).

The ACD Model and FCNC Processes

As our analysis of [241, 83] shows, the ACD model with one extra dimension has a number of interesting properties from the point of view of FCNC processes discussed here. These are:

- The GIM mechanism [246] that significantly improves the convergence of the sum over the KK modes corresponding to the top quark, removing simultaneously to an excellent accuracy the contributions of the KK modes corresponding to lighter quarks and leptons. This feature removes the sensitivity of the calculated branching ratios to the scale $M_s \gg 1/R$ at which the higher dimensional theory becomes non-perturbative, and at which the towers of the KK particles must be cut off in an appropriate way. This should be contrasted with models with fermions localized on the brane, in which the KK parity is not conserved, and the sum over the KK modes diverges. In these models the results are sensitive to m_s and, for instance, in $\Delta m_{s,d}$, the KK effects are significantly larger [247] than found by us. We expect similar behavior in other processes considered below.
- The low energy effective Hamiltonians are governed by local operators already present in the Standard Model. As flavor violation and CP violation in this model is entirely governed by the CKM matrix, the ACD model belongs to the class of models with minimal flavor violation (MFV), as defined in [78]. This has automatically the following important consequence for the FCNC processes considered in [241, 83]: the impact of the KK modes on the processes in question amounts only to the modification of the Inami-Lim one-loop functions [248].
- Thus in the case of $\Delta m_{d,s}$ and of the parameter ε_K, that are relevant for the standard analysis of the unitarity triangle, these modifications have to be made in the function S [249]. In the case of the rare K and B decays that are dominated by Z^0 penguins the functions X and Y [250] receive KK contributions. Finally, in the case of the decays $B \to X_s\gamma$, $B \to X_s$ gluon, $B \to X_s\mu\overline{\mu}$ and $K_S^0 \to \pi^0 e^+ e^-$ and the CP-violating ratio ε'/ε the KK contributions to new short distance functions have to be computed. These are the functions D (the γ penguins), E (gluon penguins), D' (γ-magnetic penguins) and E' (chromomagnetic penguins). Here we will only report on the decays relevant for Super B Factories.

Thus, each function mentioned above, which in the Standard Model depends only on m_t, now also becomes a function of $1/R$:

$$F(x_t, 1/R) = F_0(x_t) + \sum_{n=1}^{\infty} F_n(x_t, x_n), \quad F = B, C, D, E, D', E', \tag{5.1}$$

with x_n defined in (5.169). The functions $F_0(x_t)$ result from the penguin and box diagrams in the Standard Model and the sum represents the KK contributions to these diagrams.

In phenomenological applications, it is convenient to work with the gauge invariant functions [250]

$$X = C + B^{\nu\overline{\nu}}, \qquad Y = C + B^{\mu\overline{\mu}}, \qquad Z = C + \frac{1}{4}D. \tag{5.2}$$

The functions $F(x_t, 1/R)$ have been calculated in [241, 83] with the results given in Table 5-17. Our results for the function S have been confirmed in [251]. For $1/R = 300$ GeV, the functions S, X, Y, Z are enhanced by 8%, 10%, 15% and 23% relative to the Standard Model values, respectively. The impact of the KK modes on the function D is negligible. The function E is moderately enhanced but this enhancement plays only a marginal role in the phenomenological applications. The most interesting are very strong suppressions of D' and E', that for $1/R = 300\,\mathrm{GeV}$ amount to 36% and 66% relative to the Standard Model values, respectively. However, the effect of the latter suppressions is softened in the relevant branching ratios through sizable additive QCD corrections.

Table 5-17. *Values for the functions S, X, Y, Z, E, D', E', C and D.*

$1/R$ [GeV]	S	X	Y	Z	E	D'	E'	C	D
200	2.813	1.826	1.281	0.990	0.342	0.113	−0.053	1.099	−0.479
250	2.664	1.731	1.185	0.893	0.327	0.191	0.019	1.003	−0.470
300	2.582	1.674	1.128	0.835	0.315	0.242	0.065	0.946	−0.468
400	2.500	1.613	1.067	0.771	0.298	0.297	0.115	0.885	−0.469
Standard Model	2.398	1.526	0.980	0.679	0.268	0.380	0.191	0.798	−0.476

The impact of the KK modes on specific decays

The impact on the Unitarity Triangle. The function S plays the crucial role here. Consequently the impact of the KK modes on the Unitarity Triangle is rather small. For $1/R = 300\,\mathrm{GeV}$, $|V_{td}|$, $\overline{\eta}$ and γ are suppressed by 4%, 5% and 5°, respectively. It will be difficult to see these effects in the $(\overline{\varrho}, \overline{\eta})$ plane. On the other hand, a 4% suppression of $|V_{td}|$ means an 8% suppression of the relevant branching ratio for rare decays sensitive to $|V_{td}|$ and this effect has to be taken into account. Similar comments apply to $\overline{\eta}$ and γ. Let us also mention that for $1/R = 300\,\mathrm{GeV}$, Δm_s is enhanced by 8%; in view of the sizable uncertainty in $\hat{B}_{B_s}\sqrt{f_{B_s}}$, this will also be difficult to see.

The impact on rare B decays. Here, the dominant KK effects enter through the function C, or equivalently, X and Y, depending on the decay considered. In Table 5-18 we show seven branching ratios as functions of $1/R$ for central values of all remaining input parameters. The hierarchy of the enhancements of branching ratios can easily be explained by inspecting the enhancements of the functions X and Y that is partially compensated by the suppression of $|V_{td}|$ in decays sensitive to this CKM matrix element, but fully effective in decays governed by $|V_{ts}|$.

Table 5-18. *Branching ratios for rare B decays in the ACD model and the Standard Model as discussed in the text.*

$1/R$	200 GeV	250 GeV	300 GeV	400 GeV	Standard Model
$Br(B \to X_s \nu\overline{\nu}) \times 10^5$	5.09	4.56	4.26	3.95	3.53
$Br(B \to X_d \nu\overline{\nu}) \times 10^6$	1.80	1.70	1.64	1.58	1.47
$Br(B_s \to \mu^+\mu^-) \times 10^9$	6.18	5.28	4.78	4.27	3.59
$Br(B_d \to \mu^+\mu^-) \times 10^{10}$	1.56	1.41	1.32	1.22	1.07

For $1/R = 300\,\mathrm{GeV}$, the following enhancements relative to the Standard Model predictions are seen: $B \to X_d \nu\overline{\nu}$ (12%), $B \to X_s \nu\overline{\nu}$ (21%), $B_d \to \mu\overline{\mu}$ (23%) and $B_s \to \mu\overline{\mu}$ (33%). These results correspond to central values of the input parameters. The uncertainties in these parameters partly cover the differences between the ACD model and the Standard Model, and it is essential to considerably reduce these uncertainties if one wants to see the effects of the KK modes in the branching ratios in question.

The Impact on $B \to X_s\gamma$ and $B \to X_s$ gluon. Due to strong suppressions of the functions D' and E' by the KK modes, the $B \to X_s\gamma$ and $B \to X_s\,\mathrm{gluon}$ decays are considerably suppressed compared to Standard Model estimates (consult [252, 253] for the most recent reviews). For $1/R = 300\,\mathrm{GeV}$, $\mathcal{B}(B \to X_s\gamma)$ is suppressed by 20%, while $\mathcal{B}(B \to X_s\,\mathrm{gluon})$ even by 40%. The phenomenological relevance of the latter suppression is unclear at present as $\mathcal{B}(B \to X_s\,\mathrm{gluon})$ suffers from large theoretical uncertainties, and its extraction from experiment is very difficult, if not impossible.

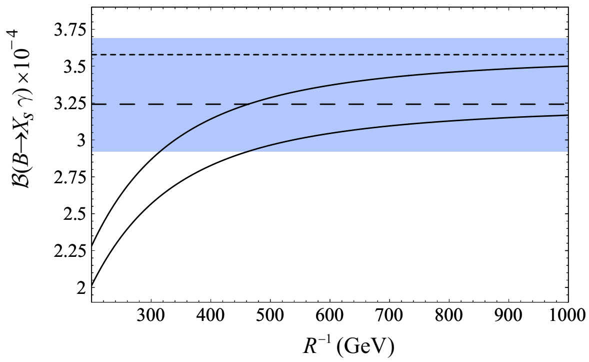

Figure 5-44. *The branching ratio for $B \to X_s\gamma$ and $E_\gamma > 1.6$ GeV as a function of $1/R$. See text for the meaning of various curves.*

In Fig. 5-44 we compare $\mathcal{B}(B \to X_s\gamma)$ in the ACD model with the experimental data and with the expectations of the Standard Model. The shaded region represents the data $\mathcal{B}(B \to X_s\gamma)_{E_\gamma>1.6\mathrm{GeV}} = (3.28^{+0.41}_{-0.36}) \cdot 10^{-4}$ [254] and the upper (lower) dashed horizontal line are the central values in the Standard Model for $m_c/m_b = 0.22$ $(m_c/m_b = 0.29)$ [255, 256]. The solid lines represent the corresponding central values in the ACD model. The theoretical errors, not shown in the plot, are roughly $\pm10\%$ for all curves.

We observe that in view of the sizable experimental error and considerable parametric uncertainties in the theoretical prediction, the strong suppression of $\mathcal{B}(B \to X_s\gamma)$ by the KK modes does not yet provide a powerful lower bound on $1/R$ and values of $1/R \geq 250\,\mathrm{GeV}$ are fully consistent with the experimental result. It should also be emphasized that $\mathcal{B}(B \to X_s\gamma)$ depends sensitively on the ratio m_c/m_b; the lower bound on $1/R$ is shifted above $400\,\mathrm{GeV}$ for $m_c/m_b = 0.29$, if other uncertainties are neglected. In order to reduce the dependence on m_c/m_b a NNLO calculation is required [255, 256, 257]. Once it is completed, and the experimental uncertainties further reduced – a Super B Factory could increase the experimental sensitivity up to a factor of three [258]–$\mathcal{B}(B \to X_s\gamma)$ may provide a very powerful bound on $1/R$ that is substantially stronger than the bounds obtained from the electroweak precision data. The suppression of $\mathcal{B}(B \to X_s\gamma)$ in the ACD model has already been found in [259]. The result presented above is consistent with the one obtained by these authors, but differs in details, as only the dominant diagrams have been taken into account in the latter paper, and the analysis was performed in the LO approximation.

The Impact on $B \to X_s\mu^+\mu^-$ and $A_{FB}(\hat{s})$. In Fig. 5-45 we show the branching ratio $\mathcal{B}(B \to X_s\mu^+\mu^-)$ as a function of $1/R$. The observed enhancement is mainly due to the function Y that enters the Wilson coefficient of the operator $(\overline{s}b)_{V-A}(\overline{\mu}\mu)_A$. The Wilson coefficient of $(\overline{s}b)_{V-A}(\overline{\mu}\mu)_V$, traditionally denoted by C_9, is essentially unaffected by the KK contributions.

Of particular interest is the forward-backward asymmetry $A_{FB}(\hat{s})$ in $B \to X_s\mu^+\mu^-$ that, similar to the case of exclusive decays [260], vanishes at a particular value $\hat{s} = \hat{s}_0$. The fact that $A_{\mathrm{FB}}(\hat{s})$ and the value of $\hat{s}_0$ being sensitive to short distance physics are in addition subject to only very small non-perturbative uncertainties makes them particularly useful quantities to test physics beyond the Standard Model . A precise measurement however is a difficult task, but it could be performed at a Super B Factory [261].

The calculations for $A_{\mathrm{FB}}(\hat{s})$ and of $\hat{s}_0$ have recently been done including NNLO corrections [262, 263] that turn out to be significant. In particular they shift the NLO value of $\hat{s}_0$ from 0.142 to 0.162 at NNLO. In Fig. 5-46 (a) we show the normalized forward-backward asymmetry that we obtained by means of the formulae and the computer program

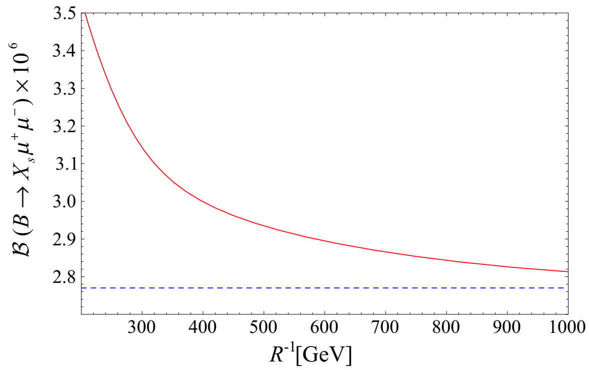

Figure 5-45. *$\mathcal{B}(B \to X_s\mu^+\mu^-)$ in the Standard Model (dashed line) [253, 52], and in the ACD model, where the dilepton mass spectrum has been integrated between the limits: $\left(\frac{2m_\mu}{m_b}\right)^2 \leq \hat{s} \leq \left(\frac{M_{J/\psi}-0.35\,\mathrm{GeV}}{m_b}\right)^2$ where $\hat{s} = (p_+ + p_-)^2/m_b^2$.*

of [52, 262] modified by the KK contributions calculated in [83]. The dependence of $\hat{s}_0$ on $1/R$ is shown in Fig. 5-46 (b).

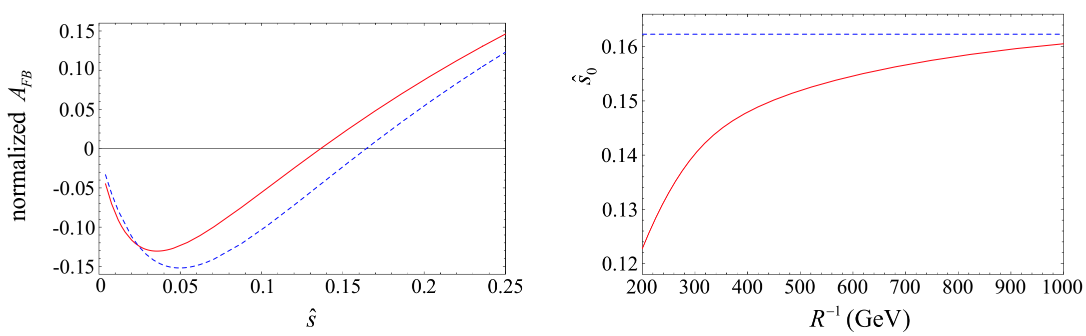

Figure 5-46. *(a) Normalized forward-backward asymmetry in the Standard Model (dashed line) and ACD for $R^{-1} = 250$ GeV. (b) Zero of the forward backward asymmetry A_{FB} in the Standard Model (dashed line) and the ACD model.*

We observe that the value of $\hat{s}_0$ is considerably reduced relative to the Standard Model result obtained by including NNLO corrections [52, 262, 263]. This decrease is related to the decrease of $\mathcal{B}(B \to X_s\gamma)$ as discussed below. For $1/R = 300\,\mathrm{GeV}$ we find the value for $\hat{s}_0$ that is very close to the NLO prediction of the Standard Model. This result demonstrates very clearly the importance of the calculations of the higher order QCD corrections, in particular in quantities like $\hat{s}_0$ that are theoretically clean. We expect that the results in Figs. 5-46 (a) and (b) will play an important role in the tests of the ACD model in the future.

In MFV models there exist a number of correlations between different measurable quantities that do not depend on specific parameters of a given model [78, 264]. In [83] a correlation between $\hat{s}_0$ and $\mathcal{B}(B \to X_s\gamma)$ has been pointed out. It is present in the ACD model and in a large class of supersymmetric models discussed for instance in [52]. We show this correlation in Fig. 5-47. We refer to [83] for further details.

Concluding Remarks

Our analysis of the ACD model shows that all the present data on FCNC processes are consistent with $1/R$ as low as $250\,\mathrm{GeV}$, implying that the KK particles could, in principle, already be found at the Tevatron. Possibly, the most

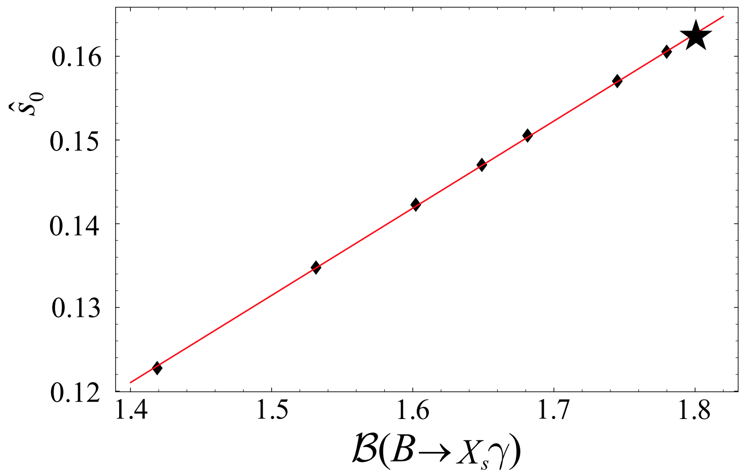

Figure 5-47. *Correlation between* $\sqrt{\mathcal{B}(B \to X_s\gamma)}$ *and* $\hat{s}_0$. *The straight line is a least square fit to a linear function. The dots are the results in the ACD model for* $1/R = 200, 250, 300, 350, 400, 600$ *and* 1000 *GeV and the star denotes the Standard Model value.*

interesting results of our analysis is the enhancement of $\mathcal{B}(K^+ \to \pi^+\nu\overline{\nu})$ (see [241] for details), the sizable downward shift of the zero ($\hat{s}_0$) in the A_{FB} asymmetry and the suppression of $\mathcal{B}(B \to X_s\gamma)$.

The nice feature of this extension of the Standard Model is the presence of only one additional parameter, the compactification scale. This feature allows a unique determination of various enhancements and suppressions relative to the Standard Model expectations. Together with a recent study [265] that found no significant difference of $S_{\phi K^0_S}$ in UED to the Standard Model prediction, we can summarize the relative deviations to the Standard Model in this model as follows

- Enhancements: $K^0_S \to \pi^0 e^+ e^-$, Δm_s, $K^+ \to \pi^+\nu\overline{\nu}$, $K_{\rm L} \to \pi^0\nu\overline{\nu}$, $B \to X_d\nu\overline{\nu}$, $B \to X_s\nu\overline{\nu}$, $K^0_S \to \mu^+\mu^-$, $B_d \to \mu^+\mu^-$, $B \to X_s\mu^+\mu^-$ and $B_s \to \mu^+\mu^-$.
- Suppressions: $B \to X_s\gamma$, $B \to X_s$ gluon, the value of $\hat{s}_0$ in the forward-backward asymmetry and ε'/ε.

We would like to emphasize that violation of this pattern in future high statistics data will exclude the ACD model. For instance, a measurement of $\hat{s}_0$ higher than the Standard Model estimate would automatically exclude this model, as there is no compactification scale for which this could bhappen. Whether these enhancements and suppressions are required by the data, or whether they exclude the ACD model with a low compactification scale, will depend on the precision of the forthcoming experiments as well as on efforts to decrease theoretical uncertainties.

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5.4.4 Warped Extra Dimensions and Flavor Violation

≻– G. Burdman –≺

Randall and Sundrum have recently proposed the use of a non-factorizable geometry in five dimensions [266] as a solution of the hierarchy problem. The metric depends on the five dimensional coordinate y and is given by

$$ds^2 = e^{-2\sigma(y)}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2\,, \qquad (5.3)$$

where x^μ are the four dimensional coordinates, $\sigma(y) = k|y|$, with $k \sim M_P$ characterizing the curvature scale. The extra dimension is compactified on an orbifold S_1/Z_2 of radius r so that the bulk is a slice of AdS_5 space between two four-dimensional boundaries. The metric on these boundaries generates two effective scales: M_P and $M_P e^{-k\pi r}$. In this way, values of r not much larger than the Planck length ($kr \simeq (11 - 12)$) can be used to generate a scale $\Lambda_r \simeq M_P e^{-k\pi r} \simeq \mathcal{O}(\mathrm{TeV})$ on one of the boundaries.

In the original RS scenario, only gravity was allowed to propagate in the bulk, with the Standard Model (Standard Model) fields confined to one of the boundaries. The inclusion of matter and gauge fields in the bulk has been extensively treated in the literature [267, 268, 269, 270, 271, 272, 273]. We are interested here in examining the situation when the Standard Model fields are allowed to propagate in the bulk. The exception is the Higgs field which must be localized on the TeV boundary in order for the W and the Z gauge bosons to get their observed masses [268]. The gauge content in the bulk may be that of the Standard Model, or it might be extended to address a variety of model building and phenomenological issues. For instance, the bulk gauge symmetries may correspond to Grand Unification scenarios, or they may be extensions of the Standard Model formulated to restore enough custodial symmetry and bring electroweak contributions in line with constraints. In addition, as was recognized in Ref. [270], it is possible to generate the fermion mass hierarchy from $\mathcal{O}(1)$ flavor breaking in the bulk masses of fermions. Since bulk fermion masses result in the localization of fermion zero-modes, lighter fermions should be localized toward the Planck brane, where their wave-function has exponentially suppressed overlap with the TeV-localized Higgs, whereas fermions with order one Yukawa couplings should be localized toward the TeV brane.

This creates an almost inevitable tension: since the lightest KK excitations of gauge bosons are localized toward the TeV brane, they tend to be strongly coupled to zero-mode fermions localized there. Thus, the flavor-breaking fermion localization leads to flavor-violating interactions of the KK gauge bosons. In particular, this is the case when one tries to obtain the correct top Yukawa coupling: the KK excitations of the various gauge bosons propagating in the bulk will have FCNC interactions with the third generation quarks. This results in interesting effects, most notably in the CP asymmetries in hadronic B decays [274].

In addition, the localization of the Higgs on the TeV brane expels the wave-function of the W and Z gauge bosons away from it resulting in a slightly non-flat profile in the bulk. This leads, for instance, to tree-level flavor changing interactions of the Z^0 [275], which in “Higgless” scenarios [276] can result in significant effects in $b \to s\ell^+\ell^-$ [277].

The KK decomposition for fermions can be written as [268, 269]

$$\Psi_{L,R}(x,y) = \frac{1}{\sqrt{2\pi r}} \sum_{n=0} \psi_n^{L,R}(x) e^{2\sigma} f_n^{L,R}(y) \, , \tag{5.4}$$

where $\psi_n^{L,R}(x)$ corresponds to the nth KK fermion excitation and is a chiral four-dimensional field. The zero mode wave functions are

$$f_0^{R,L}(y) = \sqrt{\frac{2k\pi r\,(1 \pm 2c_{R,L})}{e^{k\pi r(1\pm 2c_{R,L})} - 1}}\; e^{\pm c_{R,L}\, k\, y} \, , \tag{5.5}$$

with $c_{R,L} \equiv M_f/k$ parametrizing the 5D bulk fermion mass in units of the inverse AdS radius k. The Z_2 orbifold projection is used, so that only one of these is actually allowed, either a left-handed or a right-handed zero mode. The Yukawa couplings of bulk fermions to the TeV brane Higgs can be written as

$$S_Y = \int d^4x\, dy \sqrt{-g}\, \frac{\lambda_{ij}^{5D}}{2\, M_5}\, \overline{\Psi}_i(x,y) \delta(y - \pi r) H(x) \Psi_j(x,y) \, , \tag{5.6}$$

where λ_{ij}^{5D} is a dimensionless parameter and M_5 is the fundamental scale or cutoff of the theory. Naive dimensional analysis tells us that we should expect $\lambda_{ij}^{5D} \lesssim 4\pi$. Thus the 4D Yukawa couplings as a function of the bulk mass parameters are

$$Y_{ij} = \left(\frac{\lambda_{ij}^{5D}\, k}{M_5}\right) \sqrt{\frac{(1/2 - c_L)}{e^{k\pi r(1-2c_L)} - 1}} \sqrt{\frac{(1/2 - c_R)}{e^{k\pi r(1-2c_R)} - 1}}\, e^{k\pi r(1-c_L-c_R)} \, . \tag{5.7}$$

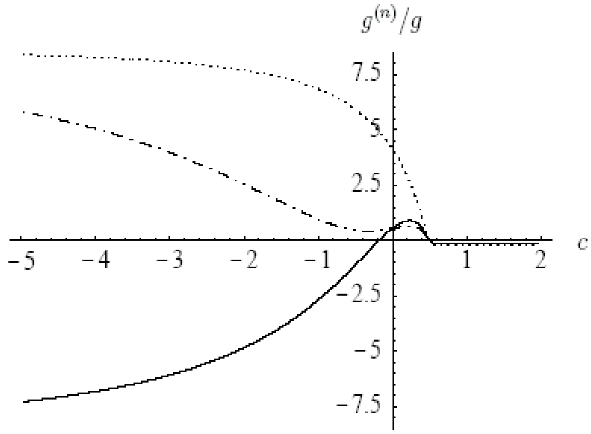


Figure 5-48. *Coupling of the first KK excitation of a gauge boson to a zero mode fermion vs. the bulk mass parameter c, normalized to the four-dimensional gauge coupling g.*

Given that we expect $k \lesssim M_5$ then the factor $\lambda_{ij}^{5D}\, k/M_5 \simeq \mathcal{O}(1)$. Thus, in order to obtain an $\mathcal{O}(1)$ Yukawa coupling, the bulk mass parameter c_L should naturally be $c_L < 0.5$ and even negative. In other words, the left-handed zero-mode should also be localized toward the TeV brane. This however, posses a problem since it means that the left-handed doublet q_L, and therefore b_L should have a rather strong coupling to the first KK excitations of gauge bosons. In Fig. 5-48 we plot the coupling of the first KK excitation of a gauge boson to a zero-mode fermion vs. the fermion's bulk mass parameter c [270].

Thus, the localization of the third generation quark doublet q_L leads to potentially large flavor violations, not only with the top quark, but also with b_L.

This induced flavor violation of KK gauge bosons with b_L (we assume b_R localized on the Planck brane) is, in principle, constrained by the precise measurement of the $Z^0 \to b\bar{b}$ interactions at the Z^0-pole. For instance, Ref. [307] considers a $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ gauge theory in the bulk. After electroweak symmetry breaking the Z^0 mixes with its KK excitations, as well as with the KK modes of a Z^0'. This generates $\delta g_L^b \lesssim \mathcal{O}(1\%) g_L^b$, compatible with current bounds, as long as $c_L \gtrsim 0.3$. Even if this is considered, it still leaves a large flavor-violating coupling of the first KK excitations to the b_L, as we can see from Fig. 5-48. More generally, for instance, in the case of strong gauge coupling [277], the effects can be even larger.

Signals in $b \to s$ and other hadronic processes

As discussed above, the flavor-changing exchange of KK gluons leads to four-fermion interactions contributing to the quark level processes $b \to d\bar{q}q$ and $b \to s\bar{q}q$, with $q = u, d, s$. We are interested in contributions that are typically of $\simeq \alpha_s$ strength due to the fact that in the product of a third generation current times a lighter quark current the enhancement in the former is (at least partially) canceled by the suppression of the latter. At low energies, the $b \to d_i\bar{q}q$ processes are described by the effective Hamiltonian [278]

$$\mathcal{H}\mathrm{eff.} = \frac{4G_F}{\sqrt{2}}\, V_{ub}V_{ui}^*\, [C_1(\mu)O_1 + C_2(\mu)O_2] - \frac{4G_F}{\sqrt{2}}\, V_{tb}V_{ti}^* \sum_{j=3}^{10} C_i(\mu)O_i + \mathrm{h.c.}, \tag{5.8}$$

where $i = d, s$ and the operator basis can be found in Ref. [278].

In the Standard Model, the operators $\{O_3 - O_6\}$ are generated from one-loop gluonic penguin diagrams, whereas operators $\{O_7 - O_{10}\}$ arise from one loop electroweak penguin diagrams. The Hamiltonian describing the $b \to s\bar{q}q$ decays is obtained by replacing V_{ts}^* for V_{td}^* in Eq. (5.8). Contributions from physics beyond the Standard Model affect the Wilson coefficients at some high energy scale. Additionally, New Physics could generate low energy interactions with the "wrong chirality" with respect to the Standard Model. This would expand the operator basis to include operators of the form $(\bar{s}_R \Gamma b_R)(\bar{q}_\lambda \Gamma q_\lambda)$, where Γ reflects the Dirac and color structure and $\lambda = L, R$.

The exchange of color-octet gauge bosons such as KK gluons of the Randall-Sundrum scenario generate flavor-violating currents with the third generation quarks. Upon diagonalization of the Yukawa matrix, this results in FCNCs at tree level due to the absence of a complete GIM cancellation. The off-diagonal elements of the left and right, up and down quark rotation matrices $U_{L,R}$ and $D_{L,R}$ determine the strength of the flavor violation. In the Standard Model, only the left-handed rotations are observable through $V_{\rm CKM} = U_L^\dagger D_L$. Here, $D^{bs}_{L,R}$, $D^{bd}_{L,R}$, $U^{tc}_{L,R}$, *etc.*, become actual observables.

The tree level flavor-changing interactions induced by the color-octet exchange are described by a new addition to the effective Hamiltonian that can, in general, be written, for $b \to s$ transitions, as

$$\delta\mathcal{H}_{\rm eff.} = \frac{4\pi\alpha_s}{M_G^2}\, D_L^{bb*}\, D_L^{bs}\, |D_L^{qq}|^2\, e^{-i\omega}\, \chi\, (\bar{s}_L \gamma_\mu T^a b_L)\, (\bar{q}_L \gamma^\mu T^a q_L) + {\rm h.c.}\ . \tag{5.9}$$

where ω is the phase relative to the Standard Model contribution; and $\chi \simeq \mathcal{O}(1)$ is a model-dependent parameter. For instance, $\chi = 1$ corresponds to the choice of $c_L \simeq 0$ that gives a coupling of the KK gauge boson about five times larger than the corresponding Standard Model value for that gauge coupling. An expression analogous to (5.9) is obtained by replacing d for s in it. This would induce effects in $b \to d$ processes.

From Eq. (5.9) we can see that the color-octet exchange generates contributions to all gluonic penguin operators. Assuming that the diagonal factors obey $|D_L^{qq}| \simeq 1$, these will have the form

$$\delta C_i = -\pi\alpha_s(M_G)\, \left(\frac{v}{M_G}\right)^2\, \left|\frac{D_L^{bs*}}{V_{tb}V_{ts}^*}\right|\, e^{-i\omega}\, f_i\, \chi\,, \tag{5.10}$$

where $f_3 = f_5 = -1/3$ and $f_4 = f_6 = 1$, and $v = 246$ GeV. This represents a shift in the Wilson coefficients at the high scale. We then must evolve the new coefficients down to $\mu = m_b$ by making use of renormalization group evolution [278]. The effects described by Eq. (5.10) are somewhat diluted in the final answer due to a large contribution from the mixing with O_2. Still, potentially large effects remain.

The phase ω in Eq. (5.9) is, in principle, a free parameter in most models and could be large. This is even true in the left-handed sector, since, in general, $V_{\rm CKM}$ comes from both the up and down quark rotation. Only if we were to argue that all of the CKM matrix comes from the down sector we could guarantee that $\omega = 0$. Furthermore, there is no such constraint in the right-handed quark sector.

We now examine what kind of effects these flavor-violating terms could produce. Their typical strength is given by α_s times a CKM-like factor coming from the $D_{L,R}$ off-diagonal elements connecting to b. The fact that the coupling is tree-level is somewhat compensated by the suppression factor $(v/M_G)^2$, for $M_G \simeq \mathcal{O}(1)$ TeV. Still, the contributions of Eq. (5.9) are typically larger than the Standard Model Wilson coefficients at the scale M_W, and, in fact, are comparable to the Wilson coefficients at the scale m_b. They could therefore significantly affect both the rates and the CP asymmetries.

The $b \to s\bar{s}s$ and $b \to s\bar{d}d$ pure penguin processes, such as $B_d \to \phi K^0_S$, $B_d \to \eta' K^0_S$ and $B_d \to \pi^0 K^0_S$, contain only small tree-level contamination of Standard Model amplitudes. These decays thus constitute a potentially clean test of the Standard Model, since their CP asymmetries are predicted to be a measurement of $\sin 2\beta_{J/\psi\, K^0_S}$, the same angle of the unitarity triangle as in the $b \to c\bar{c}s$ tree level processes such as $B_d \to J/\psi\, K^0_S$, up to small corrections. In order to estimate these effects and compare them to the current experimental information on these decay modes, we will compute the matrix elements of $\mathcal{H}_{\rm eff.}$ in the factorization approximation [279] as described in Ref. [280]. Although the predictions for the branching ratios suffer from significant uncertainties, we expect that these largely cancel when considering the effects in the CP asymmetries. Thus the CP asymmetries in non-leptonic $b \to s$ penguin-dominated processes constitute a suitable set of observables to test the effects of these color-octet states.

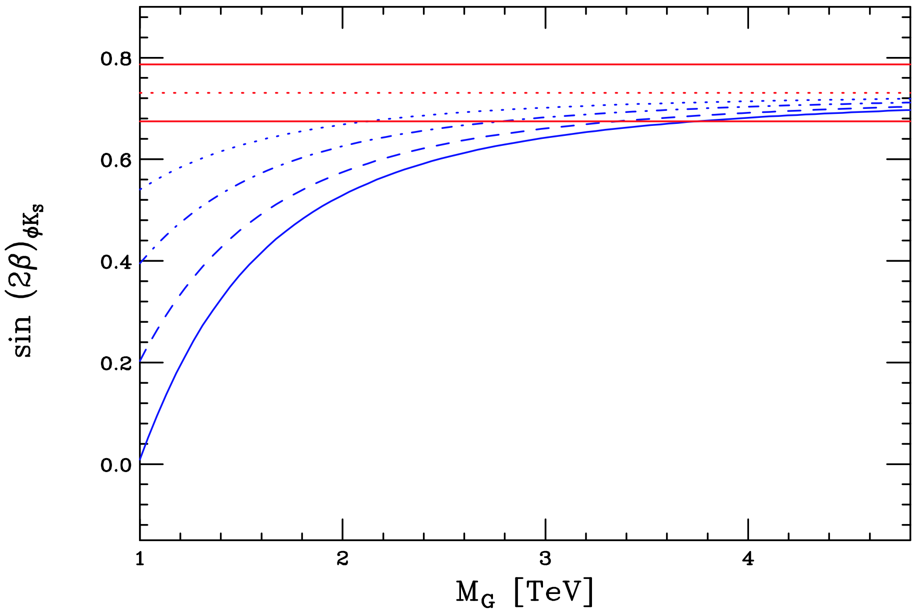


Figure 5-49. *The quantity to be extracted from the CP violation asymmetry in $B^0_d \to \phi K^0_S$ vs. the heavy gluon mass and for various values of the decay amplitude phase ω. The curves correspond to $\pi/3$ (solid), $\pi/4$ (dashed) and $\pi/6$ (dot-dash), and $\pi/10$ (dotted). The horizontal band corresponds to the world average value [281] as extracted from $B_d \to J/\psi K^0_S$, $sin(2\beta)_{\psi K^0_S} = 0.731 \pm 0.056$. From Ref. [274].*

In Fig. 5-49 we plot $\sin 2\beta_{\phi K^0_S}$ *vs.* the KK gluon mass for various values of the phase ω. Here, for concreteness, we have taken $|D^{bs}_L| = |V^*_{tb}V_{ts}|$, assumed b_R is localized on the Planck brane, and $\chi = 1$ in order to illustrate the size of the effect. The horizontal band corresponds to the $B_d \to J/\psi\, K^0_S$ measurement, $\sin 2\beta_{J/\psi\, K^0_S} = 0.731 \pm 0.056$ [281]. Only positive values of ω are shown, as negative values increase $\sin 2\beta$, contrary to the trend in the data. We see that there are sizable deviations from the Standard Model expectation for values in the region of interest $M_G \gtrsim 1$ TeV. This will be the case as long as $|D^{bs}_L| \simeq |V_{ts}|$, and $\chi \simeq \mathcal{O}(1)$, both natural assumptions.

For D^{bs}_L, this is valid as long as a significant fraction of the corresponding CKM elements comes from the down quark rotation. On the other hand, $\chi \simeq \mathcal{O}(1)$ in all the models considered here. In addition, we have not considered the effects of of D^{bs}_R, which could make the effects even larger.

Similar effects are present in $B_d \to \eta' K^0_S$, and $B_d \to K^+K^-K^0_S$ also dominated by the $b \to s\bar{s}s$ penguin contribution; as well as in the $b \to s\bar{d}d$ mode $B_d \to \pi^0 K^0_S$ [274].

The flavor-violating exchange of the KK gluon also induces an extremely large contribution to B_s–$\overline{B}_s$ mixing, roughly given by

$$\Delta m_{B_s} \simeq 200\text{ps}^{-1} \left(\frac{|D^{bs}_L|}{\lambda^2}\right)^2 \left(\frac{2\text{ TeV}}{M_G}\right)^2 \left(\frac{g_{10}}{5}\right)^2 , \tag{5.11}$$

where $\lambda \simeq 0.22$ is the Cabibbo angle, and $g_{10} \equiv g_1/g$ represents the enhancement of the zero-mode fermion coupling to the first KK gluon with respect to the four-dimensional gauge coupling, as plotted in Fig. 5-48. The contribution of Eq. (5.11) by itself is about 10 times larger than the Standard Model one for this natural choice of parameters, and would produce B_s oscillations too rapid for observation at the Tevatron or in similar experiments.

There are also similar contributions to Δm_{B_d}, when D^{bs}_L is replaced by D^{bd}_L. These were examined in Ref. [282] in the context of topcolor assisted technicolor, a much more constrained brand of topcolor than the one we consider here. The bounds found in Ref. [282] can be accommodated, as long as $|D^{bd}_L| \lesssim |V_{td}|$, which is not a very strong constraint.

Thus, we see that the flavor violation effects of the first KK gluon excitation in Randall-Sundrum scenarios where the SU(3)$_c$ fields propagate in the bulk can be significant in non-leptonic B decays, specifically in their CP asymmetries. The dominance of these effects over those induced by "weak" KK excitations, such as KK Z^0 and Z^0's, due to the larger coupling, would explain the absence of any effects in $b \to s\ell^+\ell^-$ processes, where up to now, the data is consistent with Standard Model expectations [283]. Deviations in the CP asymmetries of $b \to s$ nonleptonic processes would naturally be the first signal of New Physics in these scenarios. These very same effects can be obtained by the exchange of the heavy gluons present in generic topcolor models.

These effects can also be obtained in generic topcolor models. It is not possible to distinguish these two sources using B physics alone. This is true of any color-octet flavor-violating gauge interaction that couples strongly to the third generation. Other model building avenues addressing fermion masses might result in similar effects. In addition, the large contributions to B_s mixing, perhaps rendering Δm_{B_s} too large to be observed, is an inescapable prediction in this scenario, as it can be seen in Eq. (5.11), but it is also present in many other New Physics scenarios that produce large effects in $b \to s$ non-leptonic decays [306].

There will also be contributions from the heavy gluons to other non-leptonic B decays, such as $B \to \pi\pi$, *etc.*These modes have less clean Standard Model predictions. However, if the deviations hinted in the current data are confirmed by data samples of $500\,\text{fb}^{-1}$, to be accumulated in the next few years, it might prove of great importance to confirm the existence of these effects in less clean modes, perhaps requiring even larger data samples. Even if the heavy gluons are directly observed at the LHC, their flavor-violating interactions will be less obvious there than from large enough B physics samples. Thus, if flavor-violating interactions are observed at the LHC, high precision B physics experiments could prove crucial to elucidate their role in fermion mass generation.

Signals in $b \to s\ell^+\ell^-$

Since the wave-function of the Z^0 is pushed away from the IR brane by the boundary conditions (or a large vev), there will be non-negligible tree-level FCNC couplings of $q^T = (t_L\ b_L)^T$ and t_R with the Z^0, since these must be localized not too far from this brane. We define the effective Zbs coupling by

$$\mathcal{L}_{Zbs} = \frac{g^2}{4\pi^2}\,\frac{g}{2\cos\theta_W}\left(Z_{bs}\,\bar{b}_L\gamma^\mu s_L + Z'_{bs}\,\bar{b}_R\gamma^\mu s_R\right)Z_\mu, \tag{5.12}$$

where Z_{bs} and Z'_{bs} encode both the one loop Standard Model as well as New Physics contributions. Up to a factor of order one, the tree-level FCNC vertex induced by the flavor-violating coupling results in [277]

$$\delta Z_{bs} \simeq -\Big(-\frac{1}{2}+\frac{1}{3}\sin^2\theta_W\Big)D_L^{bs}\,\frac{8\pi^2}{g^2}\left(\frac{v^2}{m_1^2}\right)\left(\frac{g_L^2}{\pi R\,g^2}\right)f \simeq \frac{D_L^{bs}}{2}\,\frac{g_L^2}{\pi R\,g^2}N, \tag{5.13}$$

where f is defined in Ref. [277], and in order to respect the bounds from $Z \to b\bar{b}$, $|f| \lesssim \mathcal{O}(1)$. With the natural assumption $D_L^{bs} \simeq V_{tb}^* V_{ts}$, and reasonably small brane couplings $g_L^2/\pi Rg^2 = O(1)$, the correction is of the same order as the Standard Model contribution to this vertex, which is [284] $Z_{bs}^{\rm SM} \simeq -0.04$ ($Z_{bs}^{\prime\rm SM} \simeq 0$). This leads to potentially observable effects in $b \to s\ell^+\ell^-$ decays, although the current experimental data, $|Z_{bs}| \lesssim 0.08$ [284], is not greatly constraining. The effects, however, could be larger in the case of strong bulk gauge couplings [277], which may require somewhat smaller mixing angles. The effect of Eq. (5.13) also contributes to hadronic modes, such as $B \to \phi K_S^0$, although there it must compete with the parametrically larger contributions from gluonic penguins.

$D^0\overline{D}^0$ mixing

Finally, the large flavor-violating coupling of the top quark, particularly t_R, may lead to a large contribution to $D^0\overline{D}^0$ mixing. This has contributions both from KK gluon and Z^0 exchanges and has the form [277]

$$\Delta m_D \simeq 4\pi\alpha_s\,\frac{\chi(c_R)}{2m_1^2}\,\frac{(U_R^{tu*}\,U_R^{tc})^2}{2m_D}\,\langle D^0|(\bar{c}_R\gamma_\mu u_R)(\bar{c}_R\gamma^\mu u_R)|\overline{D}^0\rangle, \tag{5.14}$$

for the KK gluon exchange. Here, U_R is the rotation matrix for right-handed up quarks, and $\chi(c_R)$ is a function of c_R which gives the enhancement due to the strong coupling of the KK gluons to t_R. For instance, for $c_R \simeq 0$ and small brane couplings, $\chi \simeq 16$. To estimate the contribution to Δm_D, we need the quark rotation matrix elements. If we take $U_R^{tu*}U_R^{tc} \simeq \sin^5\theta_C$, with $\sin\theta_C \simeq 0.2$ the Cabibbo angle, then the current experimental limit on Δm_D translates[5] into $m_1 \gtrsim 2$ TeV.In the strong bulk coupling case, $\chi(c_R)$ can be enhanced and somewhat larger c_R or smaller mixing angles may be required. The contribution from the Z^0 is generically the same order, but somewhat smaller. We thus find that the effect can be consistent with, but naturally close to, the current experimental limit. Similar contributions come from the interactions of t_L, but they are typically smaller than those from t_R, because of larger values of c.

[5]Unlike for U_L and D_L, there is, in principle, no reason why U_R must have such scaling with the Cabibbo angle.

5.4.5 Warped Extra Dimensions Signatures in B Decays

≻— K. Agashe —≺

Introduction

This section is based on [285], where the reader is referred for further details and for references.

Consider the Randall-Sundrum (RS1) model which is a compact slice of AdS$_5$,

$$ds^2 = e^{-2k|\theta|r_c}\eta^{\mu\nu}dx_\mu dx_\nu + r_c^2 d\theta^2, \; -\pi \leq \theta \leq \pi, \tag{5.15}$$

where the extra-dimensional interval is realized as an orbifolded circle of radius r_c. The two orbifold fixed points, $\theta = 0, \pi$, correspond to the "UV" (or "Planck) and "IR" (or "TeV") branes respectively. In warped spacetimes the relationship between 5D mass scales and 4D mass scales (in an effective 4D description) depends on location in the extra dimension through the warp factor, $e^{-k|\theta|r_c}$. This allows large 4D mass hierarchies to naturally arise without large hierarchies in the defining 5D theory, whose mass parameters are taken to be of order the observed Planck scale, $M_{Pl} \sim 10^{18}$ GeV. For example, the 4D massless graviton mode is localized near the UV brane, while Higgs physics is taken to be localized on the IR brane. In the 4D effective theory one then finds

$$\text{Weak Scale} \sim M_{Planck} e^{-k\pi r_c}. \tag{5.16}$$

A modestly large radius, *i.e.*, $k\pi r_c \sim \log\left(M_{Planck}/\text{TeV}\right) \sim 30$, can then accommodate a TeV-size weak scale. Kaluza-Klein (KK) graviton resonances have $\sim ke^{-k\pi r_c}$, *i.e.*, TeV-scale masses since their wave functions are also localized near the IR brane.

In the original RS1 model, it was assumed that the entire Standard Model (*i.e.*, including gauge and fermion fields) is localized on TeV brane. Thus, the effective UV cut-off for gauge and fermion fields and hence the scale suppressing higher-dimensional operators, is $\sim$ TeV, *e.g.*, the same as for Higgs. However, bounds from electroweak (EW) precision data on this cut-off are $\sim 5 - 10$ TeV, whereas those from flavor-changing neutral currents (FCNC's) (for example, $K^0\overline{K}^0$ mixing) are ~ 1000 TeV. Thus, to stabilize the electroweak scale requires fine-tuning, *e.g.*, even though RS1 explains the big hierarchy between Planck and electroweak scale, it has a "little" hierarchy problem.

Bulk fermions

A solution to this problem is to move Standard Model gauge and fermion fields into the bulk. Let us begin with how bulk fermions enable us to evade flavor constraints. The localization of the wavefunction of the massless chiral fermion mode is controlled by the c-parameter. In the warped scenario, for $c > 1/2$ ($c < 1/2$) the zero mode is localized near the Planck (TeV) brane, whereas for $c = 1/2$, the wave function is *flat*.

Therefore, we choose $c > 1/2$ for light fermions, so that the effective UV cut-off is $\gg$ TeV, and thus FCNC's are suppressed. Also this naturally results in a small $4D$ Yukawa coupling to the Higgs on TeV brane without any hierarchies in the fundamental $5D$ Yukawa. Similarly, we choose $c \ll 1/2$ for the top quark to obtain an $\mathcal{O}(1)$ Yukawa. If *left*-handed top is near TeV brane, then there are FCNC's involving b_L, as follows.

Since fermions are in the bulk, we also have $5D$ gauge fields and we can show that in this set-up *high*-scale unification can be accommodated.

Couplings of fermion to gauge KK mode

The flavor violation involving b_L is due to KK modes of gauge fields, so that we need to consider couplings of these modes to fermions. We can show that wave functions of gauge KK modes are peaked near TeV brane (just like for graviton KK modes) so that their coupling to TeV brane fields (for example, the Higgs) is enhanced compared to that of zero mode (which has a flat profile) by $\approx \sqrt{2k\pi r_c}$. Thus, the coupling of gauge KK modes to zero mode fermions, denoted by $g^{(n)}$, in terms of $g^{(0)}$ (the coupling of zero mode of gauge field) has the form:

$$\begin{aligned} g^{(n)} &\sim g^{(0)} \times \sqrt{k\pi r_c},\ c \ll 1/2 \text{ (as for the Higgs)} \\ &= 0,\ c = 1/2 \text{ (fermion profile is flat)} \\ &\sim \frac{g^{(0)}}{\sqrt{k\pi r_c}},\ c \gtrsim 1/2 \text{ (independent of } c). \end{aligned} \tag{5.17}$$

Due to this coupling, there is a shift in coupling of fermions to the physical Z^0 from integrating out gauge KK modes (see Fig. 2 in Ref. [286]):

$$\delta g^{phys} \approx \sum_n g^{(n)} (-1)^n \sqrt{2k\pi r_c} M_Z^2 / m_{KK}^{(n)\,2}. \tag{5.18}$$

This shift is universal for light fermions, since light fermions have $c > 1/2$ and thus can absorbed into the S parameter.

Choice of b_L localization

It is clear that we prefer c for $(t,b)_L \ll 1/2$ inn order to obtain a top Yukawa of ~ 1 without a too large $5D$ Yukawa, but this implies a large shift in the coupling of b_L to Z^0 (relative to that for light fermions). Thus, there is a tension between obtaining top Yukawa and not shifting the coupling of b_L to Z^0. As a compromise, we choose c for $(t,b)_L$ $\sim 0.4 - 0.3$ (corresponding to a coupling of b_L to gauge KK modes $g^{(n)}/g^{(0)} \sim \mathcal{O}(1)$: see Eq. (5.17)) so that with KK masses $\sim 3 - 4$ TeV, the shift in coupling of b_L to the Z^0 is $\sim 1\%$ (see Eq. (5.18)) which is allowed by precision electroweak data.

In order to obtain a top Yukawa ~ 1, we choose c for $t_R \ll 1/2$ and c for $b_R > 1/2$ to obtain $m_b \ll m_t$. We can further show that $3 - 4$ TeV KK masses are consistent with electroweak data (S and T parameters) provided we gauge SU(2)$_R$ in the bulk [286].

Flavor violation from gauge KK modes

The flavor-violating couplings of zero-mode fermions to gauge KK modes are a result of going from a weak/gauge to a mass eigenstate basis:

$$D_L^\dagger \text{diag}\left[g^{(n)}\left(c_{L\ d}\right), g^{(n)}\left(c_{L\ s}\right), g^{(n)}\left(c_{L\ b}\right)\right] D_L, \tag{5.19}$$

where D_L is the unitary transformation from weak to mass eigenstate basis for left-handed down quarks.

Tree level KK gluon exchange contributes to $B^0\overline{B}^0$ mixing: the coefficient of $\left(\bar{b}_L \gamma^\mu d_L\right)^2$ is

$$\left[\,(D_L)_{13}\,\right]^2 \sum_n g^{(n)\,2} / m_{KK}^{(n)\,2},$$

whereas the Standard Model box diagram contribution has coefficient

$$\sim \left(V_{tb}^* V_{td}\right)^2 g^4 / \left(16\pi^2\right) \times 1/m_W^2 \sim \left(V_{tb}^* V_{td}\right)^2 / (4\ \text{TeV})^2 .$$

Since c for $(t,b)_L \sim 0.3 - 0.4$, *e.g.*, $g^{(n)}/g^{(0)} \sim \mathcal{O}(1)$, the KK gluon exchange contribution to $B^0\overline{B}^0$ mixing is comparable to the Standard Model box diagram for $m_{KK} \sim 3 - 4$ TeV. Such a large contribution is allowed, since tree level measurements of CKM matrix elements, combined with unitarity, do not really constrain V_{td} in the Standard Model, so that the Standard Model contribution to $B^0\overline{B}^0$ mixing has a large uncertainty. Explicitly, the Standard Model contribution occurs at loop level with $g/m_W \sim v$ suppression, whereas the KK gluon contribution to $B^0\overline{B}^0$ mixing is at the tree level (with $\mathcal{O}(1)$ coupling of b_L to the gauge KK mode), but suppressed by $3 - 4$ TeV $\sim 4\pi v$ KK masses, so that the two contributions are of the same size.

In other words, due to the large top mass, c for b_L is smaller (coupling to the gauge KK modes, $g^{(n)}$, is larger) than expected from m_b. This induces a large deviation from universality of the KK gluon coupling to left-handed down-type quarks. This is similar to the Standard Model, where there is no GIM suppression (due to the large m_t) in $b \to s, d$ or in the imaginary part of $s \to d$ (as opposed to the real part of $s \to d$)

This also shows that for c for $(t,b)_L \lesssim 0.3$ (*e.g.*, coupling of KK gluon to $b_L \gtrsim O(1)$), $B^0\overline{B}^0$ mixing requires $m_{KK} \gtrsim 4$ TeV—this is an independent (of $Z \to b\bar{b}$) lower limit on c for $(t,b)_L$, given that $m_{KK} \lesssim 4$ TeV by naturalness.

We next consider $b \to s\bar{s}s$. The KK gluon coupling to s is suppressed by $\sim 1/\sqrt{k\pi r_c}$. Hence we see that the effect of tree level exchange of a KK gluon in this decay is smaller than the Standard Model QCD penguin, for a choice of parameters for which the effect in $B^0\overline{B}^0$ mixing is comparable to the Standard Model value.

Flavor-violating coupling to the Z^0: $b \to s\ell^+\ell^-$

The direct effect of KK Z exchange is suppressed (compared to KK gluon exchange) by $\sim g_Z^2/g_s^2$. However, there is an indirect effect of KK Z^0 exchange: a shift in the coupling of b_L to *physical* Z^0 by $\sim 1\%$. In turn, this results in a flavor-violating coupling to the Z^0 after going to a mass eigenstate basis:

$$D_L^\dagger \text{diag}\left[\delta\left(g_Z^{d_L}\right), \delta\left(g_Z^{s_L}\right), \delta\left(g_Z^{b_L}\right)\right] D_L. \tag{5.20}$$

Since we have to allow $\delta\left(g_Z^{b_L}\right) \sim 1\%$, we get (relative to the the standard coupling of d_L to the Z^0)

$$b_L s_L Z \sim 1\,\%\, V_{ts}\,. \tag{5.21}$$

Using these couplings, we see that there are contributions to $b \to sf\overline{f}$ that are comparable to the Standard Model Z^0 penguin, with a coefficient $\sim V_{ts}\; g^2/\left(16\pi^2\right) g_Z^2/m_Z^2$, (roughly 1% in Eq. (5.21) comparable to the loop factor in the Standard Model Z^0 penguin)

This leads to a smoking gun signal in $b \to s\ell^+\ell^-$: the error in the theory prediction is $\sim 15\%$ (since V_{ts} in the Standard Model is constrained by tree level measurements of CKM matrix elements and unitarity, unlike V_{td}), so that the $\mathcal{O}(1)$ effect (relative to the Standard Model) is observable (current experiment error on measurement of this branching ratio is $\sim 30\%$). What is interesting is that the coupling of charged leptons to the Z^0 is almost axial, whereas the coupling to photons is vector; the coefficient of the operator with only axial coupling of leptons gets a new contribution, so that the angular distributions (forward-backward asymmetry) and spectrum of $\ell^+\ell^-$ are affected.

In $b \to s\bar{s}s$, the contribution of the Standard Model QCD penguin is larger than the Standard Model Z^0 penguin (roughly by $\sim g_Z^2/g_s^2$) and so the effect of $b_L s_L Z$ coupling (which is comparable to the Standard Model Z^0 penguin) is less than $\mathcal{O}(1)$ (roughly 20%). This might be observable.

Conclusions

To summarize, we have shown that bulk fermion profiles in RS1 can explain the hierarchies of fermion masses. With only first and second generations, FCNC's are small. However, including the third generation produces interesting effects. There is tension between obtaining a large top mass and not affecting the coupling of b_L to the Z^0: as a result, we have to compromise, and allow a shift in the coupling of b_L to Z^0 by $\sim 1\%$. This, in turn, leads to a flavor-violating coupling to the Z^0, and a smoking gun signal in $b \to s\ell^+\ell^-$. Finally, using the AdS/CFT correspondence, this RS1 model is dual to a $4D$ composite Higgs model; thus a strongly interacting Higgs sector can address flavor issues.

5.5 Lepton Flavor Violation

5.5.1 Lepton flavor-violating decays of the τ at a Super B Factory

>– O. Igonkina –<

Motivation

The lepton flavor-violating decays of the τ (LFV) are an excellent base for testing modern theoretical models such as supersymmetry, technicolor or models with extra dimensions.

Recent results from the neutrino oscillations experiments [287],[288],[289],[290] suggest that LFV decays do occur. However, the branching ratios expected in the charged lepton decays in the Standard Model with neutrino mixing alone is not more than 10^{-14}[291], while many other theories predict values of the order of $10^{-10} - 10^{-7}$, which should be within a reach of the Super B Factory. The predictions are summarized in Table 5-19. Among them are SUSY with different types of symmetry breaking, models with additional heavy neutrinos and models with extra gauge boson Z^0'.

Different LFV decays such as $\tau \to \ell\gamma$, $\tau \to \ell\ell\ell$, $\tau \to \ell hh$ (where ℓ is e or μ, and h is a hadron) have different importance for these models. Therefore, by studying each of these channels, one can discriminate between the models ,and extract or restrict their parameters.

Table 5-19. *Predictions for the branching ratios of $\tau \to \ell\gamma$ and $\tau \to \ell\ell\ell$ in different models.*

Model	$\tau \to \ell\gamma$	$\tau \to \ell\ell\ell$	Ref.
SM with lepton CKM	10^{-40}	10^{-14}	[291]
SM with left-handed heavy Dirac neutrino	$< 10^{-18}$	$< 10^{-18}$	[308]
SM with right-handed heavy Majorana neutrino	$< 10^{-9}$	$< 10^{-10}$	[309]
SM with left- and right-handed neutral singlets	10^{-8}	10^{-9}	[309]
MSSM with right-handed heavy Majorana neutrino	10^{-10}	10^{-9}	[310]
MSSM with seesaw	10^{-7}		[311]
left-right SUSY	10^{-10}	10^{-10}	[310]
SUSY SO(10)	10^{-8}		[193]
SUSY-GUT	10^{-8}		[312]
SUSY with neutral Higgs	10^{-10}	$10^{-10} - 10^{-7}$	[313],[314],[293]
SUSY with Higgs triplet		10^{-7}	[315]
gauge mediated SUSY breaking	10^{-8}		[316]
MSSM with universal soft SUSY breaking	10^{-7}	10^{-9}	[317]
MSSM with non-universal soft SUSY breaking	10^{-10}	10^{-6}	[318]
Non universal Z' (technicolor)	10^{-9}	10^{-8}	[319]
two Higgs doublet III	10^{-15}	10^{-17}	[320]
seesaw with extra dimensions	10^{-11}		[321]

The experimental situation

No signature for τ LFV decays has been yet found. The strictest upper limits, of order $10^{-7} - 10^{-6}$ (see Table 5-20), are limited by the size of the accumulated data samples; a data sample of 0.5 or 10 ab^{-1} will significantly improve our understanding of the mechanism of lepton flavor violation. We present the prospects for measuring τ LFV decays at the future Super B Factory, assuming that the center-of-mass energy and the detector performance similar to *BABAR*. The reconstruction of $\tau \to \ell\ell\ell$ and $\tau \to \ell\gamma$ is based on the unique topology of the $\tau\tau$ events at $\sqrt{s} \sim 10$ GeV, where τ's have a significant boost, and the decay products are easily separated. The former decays are selected using a 1-3 topology, while the latter satisfy a 1-1 topology. Additional requirements are: a positive identification of the leptons

Table 5-20. *The current strictest upper limits on the branching ratios of $\tau \to \ell\gamma$ and $\tau \to \ell\ell\ell$. The $\bar{b}$ $\tau \to \ell\ell\ell$results were published while these Proceedings were in preparation. The Belle $\tau \to \mu\gamma$, result following [292], is shown.*

$\mathcal{B}(\tau \to \ell\gamma)$	$< 3 \cdot 10^{-6}$	CLEO	(4.8 fb^{-1})	[322], [323]
$\mathcal{B}(\tau \to \mu\gamma)$	$< 2 \cdot 10^{-6}$	*BABAR*(preliminary)	(63 fb^{-1})	[324]
$\mathcal{B}(\tau \to \mu\gamma)$	$< 5 \cdot 10^{-7}$	Belle	(86.3 fb^{-1})	[325]
$\mathcal{B}(\tau \to \ell\ell\ell)$	$< 2 \cdot 10^{-6}$	CLEO	(4.8 fb^{-1})	[326]
$\mathcal{B}(\tau \to \ell\ell\ell)$	$< 3 \cdot 10^{-7}$	Belle (preliminary)	(48.6 fb^{-1})	[327]
$\mathcal{B}(\tau \to \ell\ell\ell)$	$< 1 - 3 \cdot 10^{-7}$	*BABAR*	(82 fb^{-1})	[328]
$\mathcal{B}(\tau \to \ell hh)$	$< 2 - 15 \cdot 10^{-6}$	CLEO	(4.8 fb^{-1})	[326]

from the signal decay and several kinematic cuts on the tracks in the event. The determination of the number of signal events (or the setting of an upper limit) is based on the reconstructed invariant mass and the energy of the candidates on the signal side.

Table 5-21. *Expected signal efficiency, expected background level and the sensitivity to upper limit on LFV τ decays at 90% CL for different sizes of data samples.*

		$\tau \to \ell\ell\ell$	
	90 fb^{-1}	0.5 ab^{-1}	10 ab^{-1}
Efficiency	8.5%	8%	7%
Background	0.4	1	1
UL Sensitivity	$2 \cdot 10^{-7}$	$4 \cdot 10^{-8}$	$3 \cdot 10^{-9}$
		$\tau \to \ell\gamma$	
	63 fb^{-1}	0.5 ab^{-1}	10 ab^{-1}
Efficiency	5%	4%	4%
Background	8	8	180
UL Sensitivity	$1 \cdot 10^{-6}$	$2 \cdot 10^{-7}$	$3 \cdot 10^{-8}$

The main backgrounds for $\tau \to \ell\ell\ell$ are hadronic events resulting from hadron misidentification. The study shows that the $\tau \to \ell\ell\ell$ decay is well-controlled by cuts. The required suppression of the background for $0.5\,\text{ab}^{-1}$ is achieved by additional kinematic cuts on the 1-prong side; strengthening lepton identification is essential for the analysis of the $10\,\text{ab}^{-1}$ sample. The decays $\tau \to \ell\gamma$ are contaminated with non-LFV process $\tau \to \ell\nu\nu\gamma$, which is more difficult to suppress. However, the gain due to the large statistics sample is still significant, in spite of high level of background. Table 5-21 shows the upper limit sensitivity if no signal is observed. The calculation of upper limits is done following [292]. Further improvement can be made if a new detector has better lepton identification (in particular for soft leptons, with momenta below 0.5 GeV), more accurate momentum reconstruction and larger acceptance. Precise reconstruction of the photon energy is a key issue for the analysis of the $\tau \to \ell\gamma$ decay.

By comparing Tables 5-19 and 5-21 one can see that a sample of 10 ab^{-1} will provide extremely interesting measurements. Such a measurement of $\tau \to \mu\gamma$ will be sensitive to a GUT scale m_0 up to 200 GeV, while observation of $\tau \to \mu\mu\mu$ will be sensitive to the slepton mass. It is interesting to notice that according to [293] the $\mathcal{B}(\tau \to \ell\ell\ell)$ is of the order of 10^{-7} if supersymmetric particles are heavier than 1 TeV. In such a scenario, the Super B Factory would play an essential complementary role to the LHC and ILC in exploring supersymmetry.

5.5.2 Lepton flavor-violating τ Decays in the Supersymmetric Seesaw Model

≻– J. Hisano and Y. Shimizu –≺

Introduction

The discovery of atmospheric neutrino oscillation by the SuperKamiokande experiment [294] showed that the lepton sector has a much different flavor structure than the quark sector [295][296]. The mixing angles between the first and second and between second and third generations of neutrinos are almost maximal, and these are different from the naive expectation in the grand unified theories. Many attempts to understand those mismatches between quark and lepton mixing angles have been made.

Lepton-flavor violation (LFV) in the charged lepton sector is an important tool to probe the origin of neutrino masses, if the Standard Model is supersymmetric. The finite but small neutrino masses do not predict accessible event rates for the charged LFV processes in experiments in the near future, since these processes are suppressed by the small neutrino masses, in other words, by the scale for the origin of the neutrino masses. However, if the SUSY breaking terms are generated at the higher energy scale than that for the origin of the neutrino masses, imprints may be generated in the SUSY breaking slepton mass terms, and may induce sizable event rates for charged LFV processes, which are not necessarily suppressed by the scale for the origin of the neutrino masses. Charged LFV is a thus window into the origin of neutrino masses and can lead to an understanding within supersymmetry of the observed large neutrino mixing angles.

In this article, we review LFV τ decays in the minimal SUSY seesaw model. The seesaw model is the most fascinating model to explain the small neutrino masses in a natural way [297]. This model should be supersymmetric, since it introduces a hierarchical structure between the Standard Model and the right-handed neutrino mass scale. Thus, charged LFV processes, including LFV τ decays, may be experimentally accessible in the near future.

In the next section, we review the relations between neutrino oscillation and charged LFV processes in the minimal SUSY seesaw model. In Section 5.5.2 we discuss the LFV τ decays in this model. Section 5.5.2 is devoted to discussion.

The Minimal SUSY Seesaw Model

We consider the minimal SUSY Standard Model with three additional heavy singlet-neutrino superfields ${N^c}_i$, constituting the minimal SUSY seesaw model. The relevant leptonic part of its superpotential is

$$W = N_i^c (Y_\nu)_{ij} L_j H_2 - E_i^c (Y_e)_{ij} L_j H_1 + \frac{1}{2} {N^c}_i \mathcal{M}_{ij} N_j^c + \mu H_2 H_1 \,, \tag{5.22}$$

where the indexes i, j run over three generations and $(M_N)_{ij}$ is the heavy singlet-neutrino mass matrix. In addition to the three charged-lepton masses, this superpotential has eighteen physical parameters, including six real mixing angles and six CP-violating phases. Nine parameters associated with the heavy-neutrino sector cannot be directly measured. The exception is the baryon number in the universe, if the leptogenesis hypothesis is correct [298].

At low energies the effective theory, after integrating out the right-handed neutrinos, is given by the effective superpotential

$$W_{\rm eff} = E_i^c (Y_e)_i L_j H_1 + \frac{1}{2v^2 \sin^2\beta} (\mathcal{M}_\nu)_{ij} (L_i H_2)(L_j H_2) \,, \tag{5.23}$$

where we work in a basis in which the charged-lepton Yukawa coupling constants are diagonal. The second term in (5.23) leads to light neutrino masses and to mixing. The explicit form of the small neutrino mass matrix $\mathcal{M}_\nu$ is given by

$$(\mathcal{M}_\nu)_{ij} = \sum_k \frac{(Y_\nu)_{ki}(Y_\nu)_{kj}}{M_{N_k}} v^2 \sin^2\beta \,. \tag{5.24}$$

The light neutrino mass matrix $\mathcal{M}_\nu$ (5.24) is symmetric, with nine parameters, including three real mixing angles and three CP-violating phases. It can be diagonalized by a unitary matrix U as

$$U^T \mathcal{M}_\nu U = \mathcal{M}_\nu \,. \tag{5.25}$$

By redefinition of fields, one can rewrite $U \equiv VP$, where $P \equiv \mathrm{diag}(e^{i\phi_1}, e^{i\phi_2}, 1)$ and V is the MNS matrix, with the three real mixing angles and the remaining CP-violating phase.

If the SUSY breaking parameters are generated above the right-handed neutrino mass scale, the renormalization effects may induce sizable LFV slepton mass terms, which lead to charged LFV processes [299]. If the SUSY breaking parameters at the GUT scale are universal, off-diagonal components in the left-handed slepton mass matrix $m^2_{\tilde{L}}$ and the trilinear slepton coupling A_e take the approximate forms

$$\begin{aligned}(\delta m^2_{\tilde{L}})_{ij} &\simeq -\frac{1}{8\pi^2}(3m_0^2 + A_0^2)H_{ij}\,,\\ (\delta A_e)_{ij} &\simeq -\frac{1}{8\pi^2}A_0 Y_{e_i} H_{ij}\,,\end{aligned} \tag{5.26}$$

where $i \neq j$, and the off-diagonal components of the right-handed slepton mass matrix are suppressed. Here, the Hermitian matrix H, whose diagonal terms are real and positive, is defined in terms of Y_ν and the heavy neutrino masses M_{N_k} by

$$H_{ij} = \sum_k (Y_\nu^\dagger)_{ki}(Y_\nu)_{kj} \log\frac{M_G}{M_{N_k}}, \tag{5.27}$$

with M_G the GUT scale. In Eq. (5.26) the parameters m_0 and A_0 are the universal scalar mass and trilinear coupling at the GUT scale. We ignore terms of higher order in Y_e, assuming that $\tan\beta$ is not extremely large. Thus, the parameters in H may, in principle, be determined by the LFV processes of charged leptons [300].

The Hermitian matrix H has nine parameters, including three phases, which are clearly independent of the parameters in $\mathcal{M}_\nu$. Thus $\mathcal{M}_\nu$ and H together provide the required eighteen parameters, including six CP-violating phases, by which we can parameterize the minimal SUSY seesaw model.

Our ability to measure three phases in the Hermitian matrix H, in addition to the Majorana phases $e^{i\phi_1}$ and $e^{i\phi_2}$, arelimited at present. Only a phase in H might be determined by T-odd asymmetries in $\tau \to \ell\ell\ell$ or $\mu \to eee$ [301], since they are proportional to a Jarlskog invariant obtainable from H,

$$J = \mathrm{Im}[H_{12}H_{23}H_{31}]\,. \tag{5.28}$$

However, the asymmetries, arising from interference between phases in H and $\mathcal{M}_\nu$, must be measured in order to determine other two phases in H. A possibility to determine them might be the EDMs of the charged leptons [302]. The threshold correction due to non-degeneracy of the right-handed neutrino masses might enhance the imaginary parts of the diagonal component in A_e, which contribute to the EDMs of the charged leptons. These depend on all the phases in $\mathcal{M}_\nu$ and H. However, detailed studies show that the electron and muon EDMs are smaller than $10^{-27}\,e$ cm and $10^{-29}\,e$ cm, respectively, in the parameter space where the charged LFV processes are suppressed below the experimental bounds [300].

LFV τ decays in the SUSY Seesaw Model

As explained in the previous section, we have shown that charged LFV processes give information about the minimal SUSY seesaw model which is independent of neutrino oscillation experiments. In this section we demonstrate this by considering LFV τ decays.

In the SUSY models, the LFV processes of the charged leptons are radiative, due to R parity. Thus, the largest LFV decay processes are $\tau \to \mu\gamma$ or $\tau \to e\gamma$, which come from diagrams such as those shown in Fig. 5-50. Other processes are suppressed by $\mathcal{O}(\alpha)$. These are discussed later.

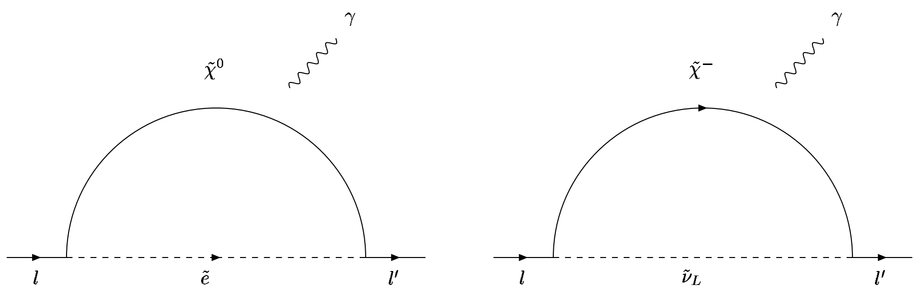


Figure 5-50. *$l \to l'\gamma$ processes in SUSY models.*

We will study LFV τ decays in two different limits of the parameter matrix H, of the form

$$H_1 = \begin{pmatrix} a & 0 & 0 \\ 0 & b & d \\ 0 & d^\dagger & c \end{pmatrix} , \tag{5.29}$$

and

$$H_2 = \begin{pmatrix} a & 0 & d \\ 0 & b & 0 \\ d^\dagger & 0 & c \end{pmatrix} , \tag{5.30}$$

where a, b, c are real and positive, and d is a complex number. The non-vanishing $(2,3)$ component in H_1 leads to $\tau \to \mu\gamma$ while the $(1,3)$ component in H_2 leads to $\tau \to e\gamma$.

In the above *ansatz*, we take $H_{12} = 0$ and $H_{13}H_{32} = 0$ because these conditions suppress $\mathcal{B}(\mu \to e\gamma)$. It is found from the numerical calculation that $\mathcal{B}(\mu \to e\gamma)$ is suppressed in a broad range of parameters with the chosen forms H_1 and H_2 [300]. From the viewpoint of model-building, the matrix H_1 is favored, since it is easier to explain the large mixing angles observed in the neutrino oscillation experiments by the structure of the Yukawa coupling Y_ν. If we adopt H_2, we might have to require some conspiracy between Y_ν and $\mathcal{M}$. However, from the viewpoint of a bottom-up approach, we can always find parameters consistent with the observed neutrino mixing angles for both H_1 and H_2, as explained in the previous section.

In Fig. 5-51 we show $\mathcal{B}(\tau \to \mu\gamma)$ for the *ansatz* H_1 and $\mathcal{B}(\tau \to e\gamma)$ for H_2 as functions of the lightest stau mass. We take the SU(2) gaugino mass to be 200 GeV, $A_0 = 0$, $\mu > 0$, and $\tan\beta = 30$ for the SUSY breaking parameters in the SUSY Standard Model. We sample the parameters in H_1 or H_2 randomly in the range $10^{-2} < a, b, c, |d| < 10$, with distributions that are flat on a logarithmic scale. Also, we require the Yukawa coupling-squared to be smaller than 4π, so that Y_ν remains perturbative up to M_G.

In order to fix $\mathcal{M}_\nu$, we fix the light neutrino parameters: $\Delta m^2_{32} = 3 \times 10^{-3}$ eV^2, $\Delta m^2_{21} = 4.5 \times 10^{-5}$ eV^2, $\tan^2\theta_{23} = 1$, $\tan^2\theta_{12} = 0.4$, $\sin\theta_{13} = 0.1$ and $\delta = \pi/2$. The Majorana phases $e^{i\phi_1}$ and $e^{i\phi_2}$ are chosen randomly. In Fig. 5-51 we assume the normal hierarchy for the light neutrino mass spectrum; as expected, the branching ratios are insensitive to the structure of the light neutrino mass matrix [300].

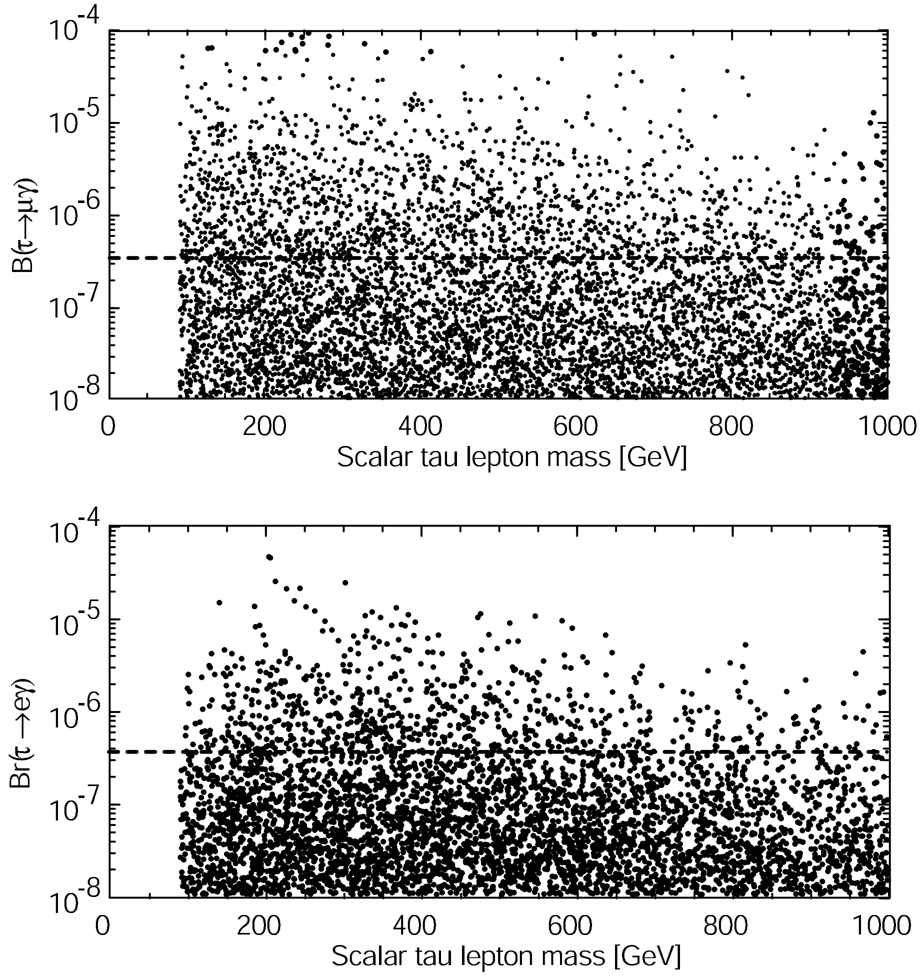


Figure 5-51. *$\mathcal{B}(\tau \to \mu\gamma)$ for H_1 and $\mathcal{B}(\tau \to e\gamma)$ for H_2. The input parameters are given in the text.*

Current experimental bounds for branching ratios of the LFV tau decays are derived by the Belle experiment, and $\mathcal{B}(\tau \to \mu(e)\gamma) < 3.2(3.6) \times 10^{-7}$ [303]. These results already exclude a fraction of the parameter space of the minimal SUSY seesaw model. These bounds can be improved to 10^{-8} at Super B Factories[303].[6]

Discussion

In this article we discussed LFV τ decays in the minimal SUSY seesaw model. We showed that the these processes provide information about the structure of this model which is independent of the neutrino oscillation experiments. Current experimental searches for $\tau \to \mu\gamma$ and $\tau \to e\gamma$ in the existing B factories already exclude a portion of the parameter space.

Let us now discuss the other LFV τ decay processes. If the SUSY breaking scale is not extremely large, $\mathcal{B}(\tau \to \mu(e)ll)$ is strongly correlated with $\mathcal{B}(\tau \to \mu(e)\gamma)$, since the on-shell photon penguin diagrams dominate

[6] If sleptons are found in future collider experiments, τ flavor violation might be found in the signal. The cross sections for $e^+e^-(\mu^+\mu^-) \to \tilde{l}^+\tilde{l}^- \to \tau^\pm\mu^\mp(\tau^\pm e^\mp) + X$ are suppressed by at most the mass difference over the widths for the sleptons. The searches for these processes in the collider experiments have more sensitivity for a small $\tan\beta$ region compared with the search for the LFV tau decays [304].

over these processes. As the result,

$$\mathcal{B}(\tau \to \mu ee(3e))/\mathcal{B}(\tau \to \mu(e)\gamma) \simeq 1/94, \quad (5.31)$$
$$\mathcal{B}(\tau \to 3\mu(e2\mu))/\mathcal{B}(\tau \to \mu(e)\gamma) \simeq 1/440, \quad (5.32)$$

where the difference between the two relations above comes from phase space. If $\tau \to \mu\gamma$ or $\tau \to e\gamma$ are found, we can perform a non-trivial test.

When the SUSY particles are very heavy, $\mathcal{B}(\tau \to \mu\gamma)$ and $\mathcal{B}(\tau \to e\gamma)$ are suppressed by the masses. However, anomalous LFV Higgs boson couplings are then generated, and these can lead to LFV processes in charged lepton decay, such as $\tau \to 3\mu$ and $\tau \to \mu\eta$ [305]. The branching ratios are limited by the muon or strange quark Yukawa coupling constant, They might, however, be observable in future experiments when $\tan\beta$ is very large and the heavier Higgs bosons are relatively light.

5.5.3 Higgs-mediated lepton flavor-violating B and τ decays in the Seesaw MSSM

≻– A. Dedes –≺

Introduction

Possible lepton number violation by two units ($\Delta L = 2$), can arise in the Standard Model or the Minimal Supersymmetric Standard Model (MSSM) from a dimension 5 operator [329, 330]

$$\Delta\mathcal{L} = -\frac{1}{4\Lambda}\, C^{AB}\,(\epsilon_{ab}\, H_a\, l_b^A)\,(\epsilon_{cd}\, H_c\, l_d^B)\, +\, \text{h.c.}\,, \quad (5.33)$$

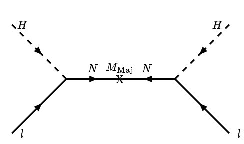
where $l^A = (\nu^A, e^A)^T$, $H = (H^+, H^0)^T$ and $A = e, \mu, \tau$ and $a, b, c, d = 1, 2$ denote lepton, Higgs doublets and SU(2) doublet indices respectively. The operator (5.33) is generated at a scale Λ, and after electroweak symmetry breaking, results in neutrino masses

$$\mathcal{L}^{\nu}_{\text{mass}} = -\frac{1}{2}\,(\frac{v^2}{4\Lambda})\, C^{AB}\, \nu^A \nu^B\, +\, \text{h.c.}\,, \quad (5.34)$$

of the order $m_\nu \simeq m_t^2/\Lambda$ and for $m_\nu = 10^{-3} - 1$ eV the scale Λ should lie in the region $10^{13} - 10^{15}$ GeV. It should be emphasized here that there is no reason for the matrix C^{AB} to be diagonal. In a basis where the charged lepton mass matrix is diagonal, C^{AB} is the infamous MNS matrix with, as atmospheric neutrino oscillation data suggest, a maximal '$23'$' or '$32'$' element. There is a mechanism which explains the existence of the operator (5.33), namely the seesaw mechanism. One can add one or more SU(2)$_L\times$ U(1)$_Y$ singlet leptonic fields N to the Lagrangian [297]

$$\Delta\mathcal{L} = -\epsilon_{ab}\, H_a\, N^A\, {Y_\nu}^{AB}\, l_b^B\, -\, \frac{1}{2}\, M^{AB}_{\text{Maj}}\, N^A N^B\, +\, \text{h.c.}\,, \quad (5.35)$$

which via the diagram



reproduces the operator (5.33) with

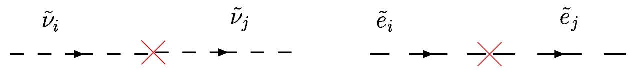
$$\Lambda^{-1} C = {Y_\nu}^T (M_{\text{Maj}})^{-1} Y_\nu\,. \quad (5.36)$$

Thus, the scale Λ is identified with the heavy Majorana mass scale M_{Maj}, which decouples at low energies, leaving only the operator (5.33), leading to neutrino masses and possibly a non-trivial mixing matrix.

In the MSSM, the influence of the renormalizable Yukawa coupling Y_ν in the slepton mass $[m_{\tilde{L}}]$ renormalisation group equations (RGEs) running from the Planck (or GUT) scale down to the scale $M_{\rm Maj}$ induces a mixing among sleptons at low energies

$$(\Delta m^2_{\tilde{L}})_{ij} \simeq -\frac{1}{8\pi^2}(3m_0^2 + A_0^2)(Y_\nu{}^\dagger \ln\frac{M_{\rm GUT}}{M_{\rm Maj}} Y_\nu) \;, \tag{5.37}$$

and thus flavor-changing insertions such as,



for sneutrinos ($\tilde{\nu}$) and charged sleptons ($\tilde{e}$) are produced. In (5.37), m_0 and A_0 are common supersymmetry breaking masses for sleptons and trilinear couplings at the Grand Unification scale (GUT), $M_{\rm GUT}$. These insertions enter in loops and result in lepton flavor-violating (LFV) processes such as $\tau \to \mu\gamma$ [see Fig.(5-52)].

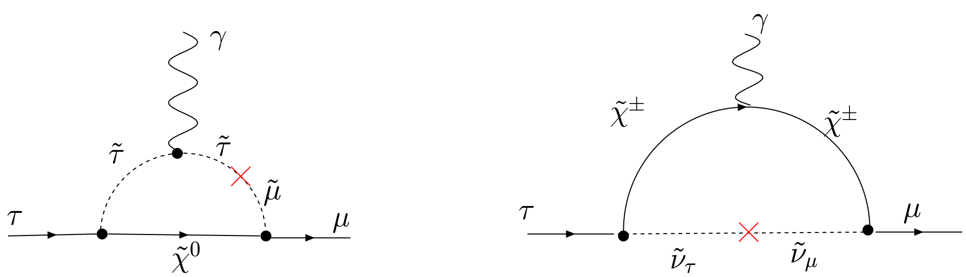


Figure 5-52. *Supersymmetric contributions to the LFV process $\tau \to \mu\gamma$.*

Equation (5.37) is the only source of lepton-flavor violation in supersymmetric seesaw models with flavor-universal soft mass terms. Since it is induced by heavy-neutrino Yukawa interactions, it relates the LFV processes to low-energy neutrino data. At the one-loop level, also gives rise to flavor violation in non-holomorphic interactions of the form $\overline{E}LH_u^*$ and leads to Higgs-mediated LFV processes in the charged-lepton sector. For technical details on deriving an effective Lagrangian for a lepton flavor-violating Higgs penguin $H - l' - l$, the reader is referred to Refs. [331, 332, 333]. Schematically, the LFV Higgs penguin is depicted in Fig. (5-53). The lepton flavor-violating Higgs penguin exists even in the Standard Model with the exchange of a W-gauge boson and the neutrino [diagram in Fig. (5-53)]. The result is proportional to the neutrino mass squared, and thus negligible. The situation changes when supersymmetrizing the $W - \nu$ loop in Fig. (5-53). Then the W-boson becomes a chargino and the neutrino becomes a sneutrino. If we calculate the chargino-sneutrino loop, we find that it is proportional to $\tan\beta$, denoted with the green dot in Fig. (5-53), and is proportional to the τ-lepton mass. In addition, the $\tau - \tau - H$ vertex is also enhanced by $\tan\beta$. Furthermore, the single neutral Higgs boson in the Standard Model becomes three neutral Higgs-bosons in the MSSM, two CP-even (h, H) and a CP-odd (A). There are, of course, other Standard Model and MSSM contributions to the penguin $\tau - \mu - H$; for simplicity these are not shown in Fig. (5-53). The last step in the derivation of the $\tau - \mu - H$ effective Lagrangian is to integrate out all the heavy particles (charginos and sneutrinos, in our case), and include the possibility of Higgs mixing (see Ref. [334] for more details), where the Higgs bosons become three CP-indistinguishable particles, H_1, H_2 and H_3, with the heaviest being the H_3. Finally, the LFV Higgs penguin and the total amplitude are enhanced by two powers of $\tan\beta$. From now on we can use the "yellow" blob (Higgs penguin) of Fig.(5-53) in the amplitude calculations for physical processes. For a complete list of applications of the Higgs penguin the reader is referred to Ref. [335]. We shall present few of them below.

Applications : $\tau \to \mu\mu\mu,\; B^0_{s,d} \to \mu\tau,\; \tau \to \mu\eta$

We shall focus on the lepton flavor-violating processes of Fig. (5-54). To quantify the above statements, we adopt, for the moment, the approximation in which we take all the supersymmetry-breaking mass parameters of the model

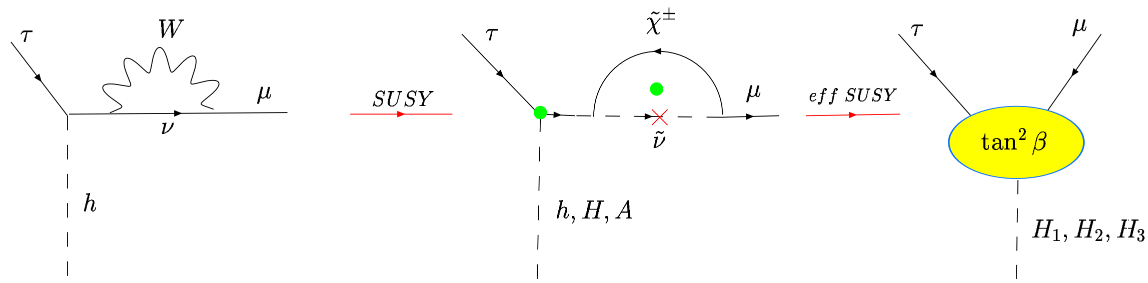


Figure 5-53. *The LFV Higgs penguin.*

to be equal at low scales, use heavy-neutrino masses that are degenerate with $M_N = 10^{14}$ GeV, and assume that $(Y_\nu^\dagger Y_\nu)_{32,33} = 1$. These are inputs for Eq. (5.37). This approximation is not realistic, but it is useful for comparing the sensitivities of different processes to New Physics: we shall present in the next section results from a more complete treatment $\tau \to 3\mu$ or related processes. In this simplified case, we obtain [332] for (5-54a)

$$Br(\tau \to 3\mu) \simeq 1.6 \times 10^{-8} \left[\frac{\tan\beta}{60}\right]^6 \left[\frac{100\,\mathrm{GeV}}{M_A}\right]^4 . \tag{5.38}$$

This should be compared with the corresponding estimate

$$Br(\tau \to \mu\gamma) \simeq 1.3 \times 10^{-3} \left[\frac{\tan\beta}{60}\right]^2 \left[\frac{100\,\mathrm{GeV}}{M_S}\right]^4 . \tag{5.39}$$

Both equations (5.38) and (5.39) are valid in the large $\tan\beta$ limit. Whereas (5.38) is two orders of magnitude *below* the present experimental bound on $\tau \to 3\mu$, (5.39) is three orders of magnitude *above* the present bound on $\tau \to \mu\gamma$. There is also the photonic penguin contribution to the decay $\tau \to 3\mu$, which is related to $\mathcal{B}(\tau \to \mu\gamma)$ by

$$\mathcal{B}(\tau \to 3\mu)_\gamma = \frac{\alpha}{3\pi}\left(\ln\frac{m_\tau^2}{m_\mu^2} - \frac{11}{4}\right)\mathcal{B}(\tau \to \mu\gamma). \tag{5.40}$$

Numerically, (5.39) leads to

$$\mathcal{B}(\tau \to 3\mu)_\gamma \simeq 3.0 \times 10^{-6} \left[\frac{\tan\beta}{60}\right]^2 \left[\frac{100\,\mathrm{GeV}}{M_S}\right]^4 , \tag{5.41}$$

which is a factor of 100 larger than (5.38). Notice also that suppressing (5.39) by postulating large slepton masses would suppress (5.38) at the same time, since sleptons enter into both loops. However, suppressing (5.38) by a large Higgs mass would not affect $\tau \to \mu\gamma$. It has recently been shown [336] that with the same assumptions, the branching ratio of the process $\tau \to \mu\eta$ [see Fig. (5-54b)], is related to that for $\tau \to 3\mu$ of Eq. (5.38)

$$\mathcal{B}(\tau \to \mu\eta)^{\mathrm{Higgs}} \;=\; 8.4 \times \mathcal{B}(\tau \to 3\mu)^{\mathrm{Higgs}} , \tag{5.42}$$

and for the double Higgs penguin diagram of Fig. (5-54c) [332]

$$\mathcal{B}(B_s^0 \to \tau\mu) \simeq 3.6 \times 10^{-7} \left[\frac{\tan\beta}{60}\right]^8 \left[\frac{100\,\mathrm{GeV}}{M_A}\right]^4 , \tag{5.43}$$

where only the leading $\tan\beta$ dependence is presented. These have to be compared with the upper limits in Table 5-22 below. In the case of B_d mesons, one should just multiply (5.43) with $|V_{td}/V_{ts}|^2 \simeq 0.05$. As expected, the Higgs-mediated branching ratio for $B_s^0 \to \tau\mu$ and $\tau \to \mu\eta$ can be larger than the one for $\tau \to 3\mu$.

Recently [337] the effect of the Higgs-exchange diagram for the lepton flavor-violating muon-electron conversion process, $\mu N \to eN$, in the supersymmetric seesaw model has been studied. The ratio of $\mathcal{B}(\mu N \to eN)/\mathcal{B}(\mu \to e\gamma)$ is enhanced at large $\tan\beta$ and for a relatively light Higgs sector. For reviews on τ-lepton and B-meson lepton flavor-violating decays, the reader should consult Refs. [338, 339].

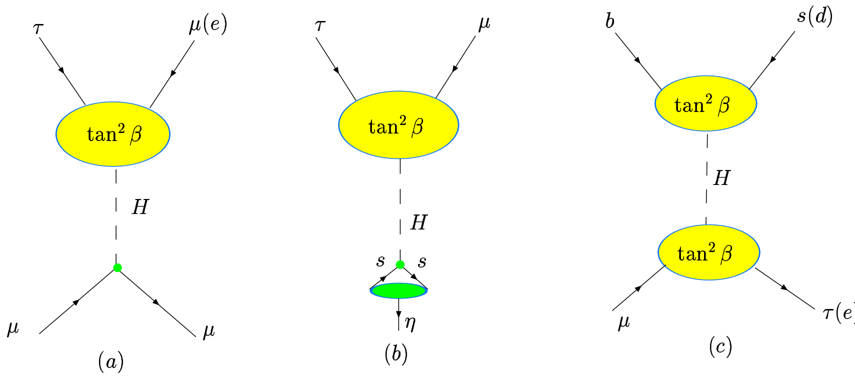


Figure 5-54. *(a) $\tau \to \mu(e)\mu\mu$ (b) $\tau \to \mu\eta$ (c) $B_{s(d)} \to \mu\tau$ Higgs-mediated processes in the MSSM. The green dot indicates an additional enhancement of $\tan\beta$.*

Results

The purpose of this work is to study in a complete way the allowed rates for the Higgs-mediated LFV processes in supersymmetric seesaw models in which the only source of LFV is the renormalization of the soft supersymmetry-breaking mass parameters [see Eq. (5.37)] above M_{N_i}, due to the singlet-neutrino Yukawa couplings of Eq. (5.35). We follow the analysis of Ref. [332] where the more general flavor-universal MSSM case was considered in which the universal masses for squarks, sleptons and the Higgs doublets H_d and H_u are different from each other. This permits different mass scales for squarks and sleptons which, in turn, are independent of the Higgs boson masses. However, we always require that squark and slepton mass matrices at the GUT scale are each proportional to unit matrices. If one goes beyond this assumption, arbitrary sources of flavor violation appear in the soft supersymmetry-breaking sector, and the model loses all predictive power, in particular the connection between the LFV and the neutrino masses and mixing. We parametrise the singlet-neutrino Yukawa couplings Y_ν and masses M_{N_i} in terms of low energy neutrino data according to Ref. [340]. We generate all the free parameters of the model randomly and calculate the low-energy sparticle masses and mixing by numerically solving the one-loop renormalization-group equations and imposing the requirement of electroweak symmetry breaking. For the decays $l_i \to l_j\gamma$, we use the exact diagrammatic formulae in [341]. The experimental CLEO and CDF bounds on $\tau \to \mu\gamma$ $B_s \to \mu\mu$ respectively, are reached first, and place upper allowed limits on $B_{s,d} \to \ell\ell'$, $\tau \to \ell\ell\ell'$ and $\tau \to \mu\eta$, that are summarized in Table 5-22.[7]

In summary, Higgs boson mediated FCNC (LFV) processes are interesting in the seesaw-MSSM. They are enhanced by powers of $\tan\beta$ and maybe orders of magnitude larger than the Standard Model predictions. By looking at Table 5-22 and comparing with the sensitivity of a Super B Factory [342] we conclude that

- $\tau \to \mu\gamma$ is (at the moment) the only LFV decay which can saturate the experimental bound. Searching for this mode at a Super B Factory is compulsory.
- LFV modes like $B \to \ell'\ell, \tau \to \ell'\ell\ell$, $\tau \to \mu\eta$, (relevant to B Factories) turn out to be small, with $\mathcal{B}$'s at $10^{-9} - 10^{-10}$, if current experimental constraints from $B_s \to \mu\mu$ and $\tau \to \mu\gamma$ are imposed.
- Searching for LFV modes like $B \to \ell'\ell, \tau \to \ell'\ell\ell$, $\tau \to \mu\eta$, could distinguish among various models for neutrino masses, for example: the R parity-violating MSSM [343] and the seesaw MSSM.

[7] After this talk was presented, newly improved bounds on $\tau \to \mu\gamma$ and $\tau \to 3\mu$ appeared [342]. These bounds set further constraints on the other processes of Table 5-22.

Table 5-22. *Experimental bounds and maximal value predictions for B- and τ decays discussed in this section. Results on $\mathcal{B} \to \ell^{+'}\ell^-$, $\tau^- \to \ell^-\ell^+\ell^-$ are based on the analysis of Ref. [332] described in the text. For $\tau \to \mu\eta$, we use the relation (5.42) from Ref. [339].*

$\mathcal{B}$(Channel)	Expt.	Bound (90% CL)	Higgs med. MSSM
$B_s \to e^+\mu^-$	CDF	$< 6.1 \times 10^{-6}$	$\lesssim 10^{-11}$
$B_s \to e^+\tau^-$	–	$--$	$\lesssim 4 \times 10^{-9}$
$B_s \to \mu^+\tau^-$	–	$--$	$\lesssim 4 \times 10^{-9}$
$B_d \to e^+\mu^-$	BABAR	$< 2.0 \times 10^{-7}$	$\lesssim 6 \times 10^{-13}$
$B_d \to e^+\tau^-$	CLEO	$< 5.3 \times 10^{-4}$	$\lesssim 2 \times 10^{-10}$
$B_d \to \mu^+\tau^-$	CLEO	$< 8.3 \times 10^{-4}$	$\lesssim 2 \times 10^{-10}$
$\tau^- \to \mu^-\mu^+\mu^-$	Belle	$< 3.8 \times 10^{-7}$	$\lesssim 4 \times 10^{-10}$
$\tau^- \to e^-\mu^+\mu^-$	Belle	$< 3.1 \times 10^{-7}$	$\lesssim 4 \times 10^{-10}$
$\tau^- \to \mu^-\eta$	CLEO	$< 9.6 \times 10^{-6}$	$\lesssim 3 \times 10^{-9}$
$\tau^- \to \mu^-\gamma$	CLEO	$< 1.1 \times 10^{-6}$	$< 1.1 \times 10^{-6}$

5.6 Conclusions and Patterns of New Physics Contributions

As discussed in the overview of this chapter, there are important questions that remain unanswered within the Standard Model. Various theoretical ideas have been proposed to answer these questions, and experimental efforts have been undertaken and will continue in the future to search for New Physics. The goal of the Super B Factory is therefore to look for New Physics effects in the Heavy Flavor sector and furthermore, to provide critical information from the measurement of new sources of CP violation and flavor mixing which will distinguish between various theoretical models.

In this chapter, we have first discussed model-independent approaches for analyzing New Physics contributions in B decays. In this case, a set of new parameters is introduced at the level of the effective Lagrangian or decay/mixing amplitude, and various observable quantities are used to constrain these parameters. Model-independent approaches have been very successful in the past, such as the determination of the Michel parameters in muon decays and the parametrization of oblique corrections in electroweak precision measurements. In B decays, depending on the processes considered, parameters representing New Physics contributions to the $B_d - \overline{B_d}$ amplitude and hadronic and electroweak $b \to s$ transitions can be introduced. In the decay $B \to V_1V_2$, where $V_{1,2}$ represent two vector particles, various angular distributions can be used to separate and measure different amplitudes, including those with CP phases. These model-independent techniques are particularly useful for a global analysis with several observables, or in comparing experimental data with the predictions of several theoretical models.

In this chapter, B physics signals in specific theoretical models are also examined. The focus is mainly on models with supersymmetry and large extra dimensions. These are leading candidates of physics beyond the Standard Model from the viewpoint of the gauge hierarchy problem. At the same time, these models have new sources of CP violation and flavor mixing structure.

- **Supersymmetry**
  Squark mass matrices contain new sources of flavor mixing and CP violation. These matrices depend on the supersymmetry breaking mechanism and flavor-dependent interactions at the high energy (GUT) scale and through renormalization the supersymmetric contributions to various B observables vary for different supersymmetry breaking scenarios. The deviations from the Standard model predictions are small in the minimal supergravity model, except for the case of $B \to D\tau\nu$, $B \to s\ell^+\ell^-$ as well as $B_s \to \mu\mu$ with possible Higgs exchange with a large ratio of the two Higgs-doublet vacuum expectation values ($\tan\beta$). The flavor mixing in the neutrino sector can be an additional source of the squark flavor mixing in the context of SUSY GUT models. In this case, large deviations from the Standard Model prediction are possible in the time-dependent CP asymmetries in the $B \to \phi K_S^0$ and $B \to K^*\gamma$ modes. This model also predicts possible observable effects in other processes such as B_s mixing, lepton flavor violation in $\tau \to \mu\gamma$, and hadronic electric dipole moments. In the scenario of effective supersymmetry where the squarks of the first two generations are very heavy, various signals with a particular pattern are expected in rare decay processes.

- **Large extra dimensions**
  Although the introduction of large extra dimensions is not directly related to flavor physics, new flavor signals are expected, particularly in the case where the observed pattern of the fermion mass spectrum is explained from some geometrical setting. The impact on B physics therefore depends on the details of the fermion mass generation mechanism. A generic signal of these theories is, however, the existence of Kaluza-Klein gravitons, and the exchange of these modes generates a set of higher dimensional operators. Such operators can produce characteristic signals in the angular distribution of the decay $b \to s\ell^+\ell^-$. The geometric construction of the fermion mass hierarchy could generate tree level flavor-changing couplings for Kaluza-Klein gauge bosons both in flat TeV^{-1} scale extra dimensions with split fermions, and in warped extra dimensions. These effects can induce large deviations from the Standard Model prediction for meson mixing as well as flavor transition diagrams. On the other hand, such deviations are not very large in the case of universal extra dimensions, where the source of flavor mixing resides solely in the CKM matrix.

Patterns of deviations from the Standard Model predictions in the various models studied in this chapter are summarized in Table 5-23. We can see that New Physics effects can appear in different processes depending on different assumptions for the origin of flavor structure in these models. It is therefore important to clarify these patterns to distinguish various models. Although comparison of various processes in B_d, B_s, K, D and lepton flavor violation are important, it is remarkable that the Super-B factory itself can provide several ways to look for New Physics contributions.

This example shows the complementary nature of flavor physics and energy frontier physics. At the LHC and ILC, direct searches for SUSY particles or Kaluza-Klein modes is essential to establish the existence of New Physics. On the other hand, there are a variety of possibilities for the origin of flavor structure within supersymmetry or models with extra dimensions. Flavor physics provides an important tool with which fundamental questions, such as how supersymmetry is broken, or how fermions propagate in extra dimensions, can be addressed, and the Super B Factory will play a central role in answering these questions.

Table 5-23. *Pattern of deviations from the Standard Model predictions in various models of supersymmetry and extra dimensions. Processes with possible large deviations are indicated. "-" means that the deviation is not expected to be large enough for observation, or not yet studied completely.*

Model	B_d Unitarity	Time-dep. CPV	Rare B decay	Other signals
mSUGRA(moderate $\tan\beta$)	-	-	-	-
mSUGRA(large $\tan\beta$)	B_d mixing	-	$B \to (D)\tau\nu$ $b \to s\ell^+\ell^-$	$B_s \to \mu\mu$ B_s mixing
SUSY GUT with ν_R	-	$B \to \phi K_S$ $B \to K^*\gamma$	-	B_s mixing τ LFV, n EDM
Effective SUSY	B_d mixing	$B \to \phi K_S$	$A_{CP}^{b\to s\gamma}$, $b \to s\ell^+\ell^-$	B_s mixing
KK graviton exchange	-	-	$b \to s\ell^+\ell^-$	-
Split fermions in large extra dimensions	B_d mixing	-	$b \to s\ell^+\ell^-$	$K^0\overline{K}^0$ mixing $D^0\overline{D}^0$ mixing
Bulk fermions in warped extra dimensions	B_d mixing	$B \to \phi K_S$	$b \to s\ell^+\ell^-$	B_s mixing $D^0\overline{D}^0$ mixing
Universal extra dimensioins	-	-	$b \to s\ell^+\ell^-$ $b \to s\gamma$	$K \to \pi\nu\overline{\nu}$

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Workshop Participants

Kaustubh Agashe
Johns Hopkins University

Justin Albert
Caltech

Ahmed Ali
DESY

David Aston
SLAC

David Atwood
University of Iowa

Henry Band
University of Wisconsin, Madison

Debbie Bard
University of Edinburgh

Christian Bauer
University of California, San Diego

Johannes Bauer
University of Mississippi

Prafulla Behera
University of Pennsylvania

Claude Bernard
Washington University, St. Louis

Jeffrey Berryhill
University of California, Santa Barbara

Adrian Bevan
University of Liverpool

Maria Biagini
INFN Laboratori Nazionali Frascati

Ikaros Bigi
University of Notre Dame du Lac

Alexander Bondar
Budker Institute of Nuclear Physics

Gerard Bonneaud
CNRS

Anders Borgland
Lawrence Berkeley National Laboratory

Sylvie Brunet
Université de Montréal

Alexander Bukin
Budker Institute of Nuclear Physics

Gustavo Burdman
Lawrence Berkeley National Laboratory

Robert Cahn
Lawrence Berkeley National Laboratory

Alessandro Calcaterra
INFN Laboratori Nazionali Frascati

Giovanni Calderini
INFN Sezione di Pisa

Massimo Carpinelli
Universita’ di Pisa and INFN

Shenjian Chen
University of Colorado

Baosen Cheng
University of Wisconsin, Madison

Farrukh Chishtie
Cornell University

Stefan Christ
University of Rostock

Jonathon Coleman
University of Liverpool

Mark Convery
SLAC

David Côté
Université de Montréal

Markus Cristinziani
SLAC

Sridhara Dasu
University of Wisconsin, Madison

Mousumi Datta
University of Wisconsin, Madison

Alakabha Datta
University of Toronto

Riccardo De Sangro
INFN Laboratori Nazionali Frascati

Athanasios Dedes
Technical University, Munich

Daniele del Re
University of California, San Diego

Durmush Demir
University of Minnesota

Francesca Di Lodovico
University of Wisconsin

Lance Dixon
SLAC

Danning Dong
SLAC

Gregory Dubois-Felsmann
Caltech

Denis Dujmic
SLAC

Eric Eckhart
Colorado State University

Ulrik Egede
Imperial College, London

Andrew Eichenbaum
University of Wisconsin, Madison

Gerald Eigen
University of Bergen

Eckhard Elsen
SLAC

Ivo Eschrich
Imperial College, London

Adam Falk
Johns Hopkins University

Thorsten Feldmann
CERN, Theory Division

Fabio Ferrarotto
INFN Sezione di Roma

Fernando Ferroni
Universita’ di Roma La Sapienza

Clive Field
SLAC

Robert Flack
Royal Holloway, Bedford College

Francesco Forti
INFN Sezione di Pisa

Mandeep Gill
Lawrence Berkeley National Laboratory

Marcello Giorgi
Universita’ di Pisa

Romulus Godang
University of Mississippi

Maarten Golterman
San Francisco State University

Stephen Gowdy
SLAC

Giacomo Graziani
CEA Saclay, DAPNIA

Benjamin Grinstein
University of California, San Diego

Andrei Gritsan
Lawrence Berkeley National Laboratory

Junji Haba
KEK

Thomas Hadig
SLAC

Paul Harrison
Queen Mary, University of London

Shoji Hashimoto
KEK

Masashi Hazumi
KEK

JoAnne Hewett
SLAC

Richard Hill
SLAC

Gudrun Hiller
University of Munich

Junji Hisano
ICRR, University of Tokyo

David Hitlin
Caltech

Mark Hodgkinson
University of Manchester

George Hou
National Taiwan University

Tobias Hurth
CERN, Theory Division

Olya Igonkina
University of Oregon

Gino Isidori
INFN Laboratori Nazionali Frascati

Yoshihito Iwasaki
KEK

Joseph Izen
University of Texas at Dallas

Alex Kagan
University of Cincinnati

Nobu Katayama
KEK

Peter Kim
SLAC

Pyungwon Ko
KAIST

Christopher Kolda
University of Notre Dame du Lac

Patrick Koppenburg
KEK

Witold Kozanecki
CEA Saclay and SLAC

Frank Krüger
Technical University, Munich

Chung Khim Lae
University of Maryland

Livio Lanceri
Universita' and INFN, Trieste

Bjorn Lange
Cornell University

Urs Langenegger
SLAC

Ben Lau
Princeton University

David Leith
SLAC

Adam Lewandowski
SLAC

Jim Libby
SLAC

Zoltan Ligeti
Lawrence Berkeley National Laboratory

Ben Lillie
SLAC

David London
Université de Montréal

John LoSecco
University of Notre Dame du Lac

Xinchou Lou
University of Texas at Dallas

Steffen Luitz
SLAC

Enrico Lunghi
University of Zurich

Vera Lüth
SLAC

Harvey Lynch
SLAC

David MacFarlane
University of California, San Diego

Giampiero Mancinelli
University of Cincinnati

Dominique Mangeol
McGill University

Aneesh Manohar
University of California, San Diego

Helmut Marsiske
SLAC

Michael Mazur
University of California, Santa Barbara

Robert Messner
SLAC

Gagan Mohanty
Queen Mary, University of London

Ajit Mohapatra
University of Wisconsin, Madison

Remi Mommsen
University of California, Irvine

Thomas Moore
University of Massachusetts

Ilya Narsky
Caltech

Nicola Neri
Universita' di Pisa and INFN

Sechul Oh
KEK

Yukiyoshi Ohnishi
KEK

Yasuhiro Okada
KEK

Alexei Onuchin
Budker Institute of Nuclear Physics

Stefano Passaggio
INFN Sezione di Genova

Popat M. Patel
McGill University

Nello Paver
Dept. of Theoretical Physics, Trieste and INFN

David Payne
University of Liverpool

Gilad Perez
Lawrence Berkeley National Laboratory

Alexey Petrov
Wayne State University

Aaron Pierce
SLAC

Maurizio Pierini
Universita' di Roma, La Sapienza

Giancarlo Piredda
INFN Roma

Steve Playfer
University of Edinburgh

Frank Porter
Caltech

Chris Potter
University of Oregon

Teela Pulliam
Ohio State University

Milind Purohit
University of South Carolina

Helen Quinn
SLAC

Matteo Rama
Universita' di Pisa

Blair Ratcliff
SLAC

Thomas Rizzo
SLAC

Douglas Roberts
University of Maryland

Steven Robertson
SLAC

Steven Robertson
McGill University

Enrico Robutti
INFN Sezione di Genova

Aaron Roodman
SLAC

Ira Rothstein
Carnegie Mellon University

Anders Ryd
Caltech

Alam Saeed
SUNY, Albany

Muhammad Saleem
SUNY, Albany

Alex Samuel
Caltech

Asish Satpathy
University of Texas at Austin

Rafe Schindler
SLAC

Alan Schwartz
University of Cincinnati

John Seeman
SLAC

Vasia Shelkov
Lawrence Berkeley National Laboratory

Yasuhiro Shimizu
Nagoya University

Maria Chiara Simani
Lawrence Livermore National Laboratory

Nita Sinha
Institute of Mathematical Sciences, Chennai

Rahul Sinha
Institute of Mathematical Sciences, Chennai

Jim Smith
University of Colorado

Arthur Snyder
SLAC

Abner Soffer
Colorado State University

Michael Sokoloff
University of Cincinnati

Evgeni Solodov
Budker Institute of Nuclear Physics

Amarjit Soni
Brookhaven National Laboratory

Iain Stewart
MIT

Dong Su
SLAC

Hiroyasu Tajima
SLAC

Makoto Tobiyama
SLAC

Walter Toki
Colorado State University

Silvano Tosi
INFN Sezione di Genova

Christos Touramanis
University of Liverpool

Georges Vasseur
CEA Saclay, DAPNIA

Jaroslav Va'vra
SLAC

Oscar Vives
University of Valencia

Steven Wagner
SLAC

Steven Weber
SLAC

Achim Weidemann
University of South Carolina

Andreas Weiler
Technical University, Munich

James Wells
University of Michigan

Guy Wilkinson
University of Oxford

David C. Williams
University of California, Santa Cruz

Stephane Willocq
University of Massachusetts

Jinwei Wu
University of Wisconsin, Madison

Daniel Wyler
University of Zurich

Masa Yamauchi
KEK

Aaron Yarritu
SLAC

Jong Yi
Iowa State University

Charles Young
SLAC

Francisco Yumiceva
University of South Carolina

Qinglin Zeng
Colorado State University


Kaluza-Klein Turris
R-Symmetrae Prevaricationis Palus
Generalis Higgisuli Lacus
Amnis Neutri cum Veriatione Saporis Antrum
Sinistrae-Dextrae Symmetriae Pons
Supergravitatis Minimae Vallis
Circa Universalis Regulae Campus