

Reevaluating Bounds on Flavor-Changing Neutral Current Parameters in R-Parity Conserving and R-Parity Violating Supersymmetry from $B^0 - \overline{B}^0$ Mixing

Jyoti Prasad Saha *

Department of Physics, Jadavpur University, Kolkata 700032, India

Anirban Kundu †‡

Universität Dortmund, Institut für Physik, D-44221 Dortmund, Germany

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Abstract

We perform a systematic reevaluation of the constraints on the flavor-changing neutral current (FCNC) parameters in R-parity conserving and R-parity violating supersymmetric models. As a typical process, we study the constraints coming from the measurements on the $B^0 - \overline{B}^0$ system on the supersymmetric δ_{13}^d parameters, as well as on the products of the λ' type R-parity violating couplings. Present data allows us to put constraints on both the real and the imaginary parts of the relevant parameters.

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*Electronic address: jyotip@juphys.ernet.in

†Electronic address: Anirban.Kundu@cern.ch

‡On leave from Department of Physics, Jadavpur University, Kolkata 700032, India. Present address: Department of Physics, University of Calcutta, 92 A.P.C. Road, Kolkata 700009, India.

1 Introduction

The answer to whether there is any new physics beyond the Standard Model (SM) is probably in the affirmative. One of the most promising candidates for new physics (NP), and also most studied, is Supersymmetry (SUSY), in both its R-parity conserving (RPC) and R-parity violating (RPV) incarnations. Unfortunately for the believers of a Theory of Everything, SUSY introduces a plethora of new particles, and even in its most constrained version, a few more arbitrary input parameters over and above to that of the SM. Thus, it has become imperative to constrain the SUSY parameter space as far as possible from existing data.

There are two aspects to this practice. First, take the experimental data and find how much space we can allow for the SUSY parameters. Second, take the bounds obtained by method 1 and see what signals one should observe in present or future experiments (and try to explain if there are any apparent anomalies in the present data). Taken together, these two methods form a strong tool to observe indirect SUSY (or for that matter, any NP) signals, which is complementary to the direct observation of the new particles in high-energy machines like the Large Hadron Collider (LHC).

Since SUSY breaks at a high scale, there is no compelling mechanism to suppress the flavor-changing neutral current (FCNC) effects once we take the renormalization group (RG) evolutions into account. The problem is most severe for the gravity-mediated SUSY breaking (SUGRA) type models. A huge number of models have been proposed to solve this problem; we will not go into them here. Rather, we will focus on an equally important area of study, *viz.*, constraining the FCNC parameters from experimental data. This shows which models are to stay and which are not.

The problem is worse in the RPV version of SUSY, since there exists a large number of FCNC couplings from the very beginning. There are no theoretical limits on these couplings except that they better not be nonperturbative even at the scale of the Grand Unified Theories (GUT). The only way to constrain the individual couplings and the products of two (or more) of them is from experimental data.

We will just discuss, as a sample process, the mixing of neutral B mesons, and the effects of RPC and RPV SUSY on it (more processes will be discussed in a subsequent publication). Why this process? The reasons are manifold: (i) The theoretical part of SM and SUSY are both well-known (including higher-order QCD corrections), apart from the uncertainties in some of the inputs; (ii) The mixing and the CP-asymmetry data have reached sufficient precision, and order-of-magnitude improvements are likely to occur in near future; (iii) The SM amplitude is one-loop, so that the RPC SUSY amplitude, which must be a one-loop process, has a fair chance of competing, even with high sparticle masses; (iv) The RPV SUSY amplitude is also one-loop and competes on the same ground (there may be even tree-level amplitudes, but the couplings are highly suppressed); (v) There are some decay channels, *e.g.*, $B \rightarrow \phi K$ [1], $B \rightarrow \eta' K$ [2], $B \rightarrow \pi\pi$ [3], which indicate that there may be signals of NP hidden in B decays (although a SM explanation is never ruled out).

Effects of RPC SUSY in $B^0 - \overline{B}^0$ mixing have been exhaustively studied in the literature, and constraints were put on the FCNC parameters of different SUSY models [4, 5, 6]. Apart from that, SUSY effects on various B decay processes (radiative, leptonic, semileptonic and nonleptonic) have been thoroughly investigated, but we are not going to discuss this aspect in the present paper [7]. A similar exercise has been performed for RPV SUSY models too [8, 9].

In this paper, we will explore the robustness of some of the FCNC parameters in RPC and RPV SUSY as quoted in the literature [5, 6, 9]. The importance of this is threefold: first, this will serve as an update of the existing results in the light of new data; second, this will show that the bounds can get substantially relaxed once we take all the SM uncertainties into account (a systematic study of the SUSY FCNC parameters, taking this point consistently into account, has not been performed as far as we know, though Ref. [6] does it in a sufficiently exhaustive way for RPC SUSY only); and third, since the constraints obtained in this paper are the most conservative ones, it should tell the experimenters what sort of signal are to be expected at the most. We will also show how robust these bounds are if there happens to be a *natural* cancellation between RPC and RPV SUSY effects.

Here one must note the limitations of the B-related experiments in constraining the NP models. Apart from the experimental uncertainties ($\sin(2\beta)$, CKM elements, branching fractions, etc.), there are a number of inherent theoretical uncertainties, most of which stem from the nonperturbative nature of low-energy QCD. To disentangle signatures of NP, one must be fairly lucky to get a sizable deviation from the SM prediction. This is precisely the reason why signatures of RPC SUSY will be unobservable in those decay modes which have a tree-level amplitude in the SM.

We will consider only the gluino-mediated box for the RPC SUSY amplitude. This constrains the mixing parameters δ_{13}^d between the first and the third generations for the down-quark sector. Though neutralino and charged-Higgs boson mediated boxes are expected to be small compared to the gluino box, an almost equally large contribution comes from the chargino diagram [10, 11]. Since the latter diagram constrains the mixing in the up-quark sector (the δ_{13}^u parameters), it will not be relevant for our future discussion. The relevant details are to be found in Section 2. For RPV SUSY we will consider only one product coupling giving rise to a new mixing amplitude to be nonzero at a time.

We will assume, just for simplicity, that NP affects only the mixing amplitude, but there is no NP in the subsequent decay processes. On the one hand this means that the CP asymmetry in the channel $B \rightarrow J/\psi K_S$ measures the phase in the box amplitude. This is easy to implement for RPV SUSY by choosing appropriate combinations of nonzero flavor-dependent couplings¹. For RPC SUSY, this assumption amounts to neglecting loop-induced SUSY contributions to tree-level SM ones, which is a safe assumption.

The QCD corrections to the box amplitude has been implemented upto Next to Leading Order (NLO), both for RPC and RPV SUSY. The relevant anomalous dimension matrix may be found in [6].

The paper is organized as follows: in Section II, we briefly recall the FCNC phenomena in RPC and RPV SUSY, and discuss our numerical inputs in Section III. Section IV deals with the constraints coming from RPC SUSY, while Section V does the same job for the RPV version. In Section VI we summarize and conclude.

¹However, this is true only for channels like $B \rightarrow J/\psi K_S$. The channel $B \rightarrow \pi^+\pi^-$, which measures α , has RPV contributions both in mixing and in subsequent decay, and they must be treated together [12].

2 $B^0 - \overline{B}^0$ Mixing in SUSY

2.1 $B^0 - \overline{B}^0$ in the SM

The off-diagonal element in the 2×2 effective Hamiltonian for the neutral B system causes the mixing between the gauge eigenstates $B^0 (\equiv \overline{b}d)$ and $\overline{B}^0 (\equiv b\overline{d})$ [13]. The mass difference between the two mass eigenstates ΔM_d is given by

$$\Delta M_d = 2|M_{12}| \quad (1)$$

where

$$\begin{aligned} M_{12} &\equiv \frac{\langle \overline{B}^0 | H_{eff} | B^0 \rangle}{2m_B} \\ &= \frac{G_F^2}{6\pi^2} (V_{td}V_{tb}^*)^2 \eta_B m_B f_B^2 B_B m_W^2 S_0(x_t), \end{aligned} \quad (2)$$

with $x_t = m_t^2/m_W^2$, and

$$S_0(x) = \frac{4x - 11x^2 + x^3}{4(1-x)^2} - \frac{3x^3 \ln x}{2(1-x)^3}. \quad (3)$$

We follow the convention of Ref. [13] for normalization of the meson wave functions. The perturbative QCD corrections are parametrized by η_B , while the nonperturbative corrections are dumped in B_B . f_B is the B meson decay constant. The subleading boxes with two charm quarks or one charm and one top quark are entirely negligible (the same holds for RPC SUSY; however, due to the nonuniversal nature of the relevant couplings, this is not true for the RPV version).

For B decays to a flavor-blind final state f (e.g., $J/\psi K_S$) where there are no nonzero CKM phases in the decay amplitude, the measured CP asymmetry is proportional to the imaginary part of the mixing amplitude. For $B^0 - \overline{B}^0$ box, this is $\sin(2\beta)$ as $\arg(V_{td}) = -\beta$ (we will implicitly assume the Wolfenstein parametrization of the CKM matrix, though the physical observables are parametrization invariant). For B_s box, the amplitude is real, so there is no CP violation in the SM (to the leading order, *i.e.*, neglecting $\mathcal{O}(\lambda^4)$ terms in the CKM matrix, where $\lambda = V_{us} \approx 0.22$).

NP adds up one (or more) new term to M_{12} . Even if it is real, the effective phase β_{eff} should change from β . Thus, ΔM_d and $A_{CP}(B \rightarrow J/\psi K_S)$, taken together, should constrain both the real and the imaginary parts of the NP amplitude.

Suppose there is one NP amplitude with a weak phase of -2ϕ so that one can write

$$M_{12} = |M_{12}^{SM}| \exp(-2i\beta) + |M_{12}^{NP}| \exp(-2i\phi). \quad (4)$$

This immediately gives the effective mixing phase β_{eff} as

$$\beta_{eff} = 0.5 \arctan \frac{|M_{12}^{SM}| \sin(2\beta) + |M_{12}^{NP}| \sin(2\phi)}{|M_{12}^{SM}| \cos(2\beta) + |M_{12}^{NP}| \cos(2\phi)} \quad (5)$$

and the mass difference between the B meson mass eigenstates as

$$\Delta M_d = 2 \left[|M_{12}^{SM}|^2 + |M_{12}^{NP}|^2 + 2|M_{12}^{SM}| |M_{12}^{NP}| \cos 2(\beta - \phi) \right]^{1/2}. \quad (6)$$

These are going to be our basic formulae. The only remaining task is to find M_{12}^{NP} .

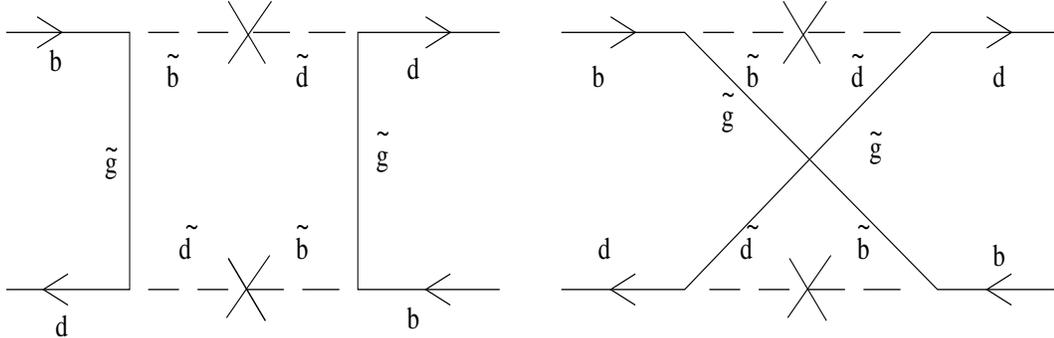


Figure 1: *Gluino-mediated SUSY contributions to $B^0 - \overline{B^0}$ mixing. One needs to add the crossed diagrams too. The crosses on the squark propagators denote the insertion of the relevant δ parameters. All chirality combinations are possible for the quarks, which we do not show explicitly.*

Note that even if $\phi = 0$, $\beta_{eff} \neq \beta$, which means that even for real NP amplitudes, the constraints vary with the choice of β . The CP asymmetry in $B \rightarrow J/\psi K_S$ measures β_{eff} . However, with NP contributing to ΔM_d , the $B^0 - \overline{B^0}$ mixing input to the standard CKM fit is lost, and β essentially becomes a free parameter. The same happens for V_{td} also. We will discuss these issues in Section III.

2.2 R-Parity conserving SUSY

In RPC SUSY models there can be two more independent phases ϕ_A and ϕ_B apart from the CKM phase, but the electric dipole moment of the neutron constrains them to be small ($\sim \mathcal{O}(10^{-2}-10^{-3})$) unless the squarks are extremely heavy or there is a fine-tuning between them [14]. We take both of them to be zero, a choice which can be theoretically motivated, since $\phi_A(\phi_B)$ is the relative phase between the common trilinear A -term (bilinear B -term) and the common gaugino mass M [15]. Even then one can have new contribution to CP violation coming from SUSY FCNC effects [4]. The origin of SUSY FCNC can be easily understood: quark and squark mass matrices are not simultaneously diagonalizable. At $q^2 \sim m_W^2$, radiative corrections induced by up-type (s)quark loops are important. These corrections are typically of the order of $\log(\Lambda_S/m_W)$ (Λ_S is the SUSY breaking scale) and hence can be large for SUGRA type models. This generates FCNC which occurs even in the quark-squark-neutral gaugino vertices, but the flavor structure is controlled by the CKM matrix. However, this last feature need not be true in any arbitrary SUSY model, particularly those with nonuniversal mass terms.

One generally works in the basis where the quark fields are eigenstates of the Hamiltonian. SUSY FCNC can be incorporated in two ways: (i) Vertex mixing, an approach where the squark propagators are flavor and ‘chirality’² conserving, and the vertices violate them; (ii) Propagator mixing, where flavor and ‘chirality’ are conserved in the vertices but changed in propagators. The second approach is more preferred for phenomenological analysis, since the higher order QCD corrections are known better. This is also known as the Mass Insertion Approximation (MIA)

²Squarks cannot have a chirality, but this is just a loose term to denote the partners of respective chiral fermions.

[16, 5].

At the weak scale one can write the 6×6 squark mass matrix (say the down type) as

$$\tilde{\mathcal{M}}_D^2 = \begin{pmatrix} \tilde{\mathcal{M}}_{DLL}^2|^{tree} + \Delta_{LL}^2 & \Delta_{LR}^2 \\ \Delta_{RL}^2 & \tilde{\mathcal{M}}_{DRR}^2|^{tree} + \Delta_{RR}^2 \end{pmatrix} \quad (7)$$

where the Δ terms incorporate the FCNC effects. Different FCNC effects are parametrized in terms of $\delta_{AB}^{ij} \equiv \Delta_{AB}^{ij}/\tilde{m}^2$, where $\tilde{m} = \sqrt{m_1 m_2}$, the geometric mean of the masses of the two participating squarks, i and j are flavor indices, and A and B are chiral indices. The δ s are completely calculable in constrained MSSM, but this is not true in general. Theoretically, for the success of perturbative analysis, one expects $|\delta| < 1$. The gluino-mediated box diagrams causing $B^0 - \bar{B}^0$ mixing constrain only δ^{13} s of different chiralities in the down-quark sector. (Thus, δ_{AB}^{13} in our notation means $(\delta_{13}^d)_{AB}$ in the more conventional notation; there is, however, no chance of confusion since we do not deal with the δ^u parameters.) The same thing can be said for the up-quark sector by considering the chargino contributions [10].

The $\Delta B = 2$ effective Hamiltonian can be written in a general form as

$$\mathcal{H}_{eff}^{\Delta B=2} = \sum_{i=1}^5 c_i O_i + \sum_{i=1}^3 \tilde{c}_i \tilde{O}_i + H.c. \quad (8)$$

where

$$\begin{aligned} O_1 &= (\bar{b}\gamma^\mu P_L d)_1 (\bar{b}\gamma_\mu P_L d)_1, \\ O_2 &= (\bar{b}P_R d)_1 (\bar{b}P_R d)_1, \\ O_3 &= (\bar{b}P_R d)_8 (\bar{b}P_R d)_8, \\ O_4 &= (\bar{b}P_L d)_1 (\bar{b}P_R d)_1, \\ O_5 &= (\bar{b}P_L d)_8 (\bar{b}P_R d)_8, \end{aligned} \quad (9)$$

and $P_{R(L)} = (1 + (-)\gamma_5)/2$. The subscripts 1 and 8 indicate whether the currents are in color-singlet or in color-octet combination. The \tilde{O}_i s are obtained from corresponding O_i s by replacing $L \leftrightarrow R$.

The Wilson coefficients have been computed at the high scale M_S (chosen to be the arithmetic mean of the average squark mass and the gluino mass) by evaluating the diagrams in Fig. 1. We quote the results [5]:

$$\begin{aligned} c_1 &= -R \left(24x f_6(x) + 66\tilde{f}_6(x) \right) (\delta_{LL}^{13})^2, \\ c_2 &= -R (204x f_6(x)) (\delta_{RL}^{13})^2, \\ c_3 &= R (36x f_6(x)) (\delta_{RL}^{13})^2, \\ c_4 &= -R \left\{ \left(504x f_6(x) - 72\tilde{f}_6(x) \right) \delta_{LL}^{13} \delta_{RR}^{13} - 132\tilde{f}_6(x) \delta_{LR}^{13} \delta_{RL}^{13} \right\}, \\ c_5 &= -R \left\{ \left(24x f_6(x) + 120\tilde{f}_6(x) \right) \delta_{LL}^{13} \delta_{RR}^{13} - 180\tilde{f}_6(x) \delta_{LR}^{13} \delta_{RL}^{13} \right\}. \end{aligned} \quad (10)$$

Here $R = \alpha_s^2/216m_q^2$ and $x = m_g^2/m_q^2$. The coefficients \tilde{c}_i s can be obtained from corresponding c_i s again with $L \leftrightarrow R$. The functions f_6 and \tilde{f}_6 are given by

$$\begin{aligned} f_6(x) &= \frac{6(1+3x)\ln x + x^3 - 9x^2 - 9x + 17}{6(x-1)^5}, \\ \tilde{f}_6(x) &= \frac{6x(1+x)\ln x - x^3 - 9x^2 + 9x + 1}{3(x-1)^5}. \end{aligned} \quad (11)$$

Next, one should evolve these coefficients down to the low-energy scale, taken, following Ref. [6], to be $\mu = m_b = 4.6$ GeV, using the NLO-QCD corrections. The low-scale Wilson coefficients are

$$c_i(\mu) = \sum_r \sum_s \left(b_r^{i,s} + \eta c_r^{i,s} \right) \eta^{ar} c_s(M_S) \quad (12)$$

where $\eta = \alpha_s(M_S)/\alpha_s(m_t)$. For the numerical values of a , b and c matrices we refer the reader to eq. (10) of Ref. [6].

The operators O_i are also to be renormalized at the scale μ . The expectation values of these operators between $\overline{B^0}$ and B^0 at the scale μ are given by

$$\begin{aligned} \langle O_1(\mu) \rangle &= \frac{2}{3} m_B^2 f_B^2 B_1(\mu), \\ \langle O_2(\mu) \rangle &= -\frac{5}{12} S m_B^2 f_B^2 B_2(\mu), \\ \langle O_3(\mu) \rangle &= \frac{1}{12} S m_B^2 f_B^2 B_3(\mu), \\ \langle O_4(\mu) \rangle &= \frac{1}{2} S m_B^2 f_B^2 B_4(\mu), \\ \langle O_5(\mu) \rangle &= \frac{1}{6} S m_B^2 f_B^2 B_5(\mu), \end{aligned} \quad (13)$$

where

$$S = \left(\frac{m_B}{m_b(m_b) + m_d(m_b)} \right)^2 \quad (14)$$

The B -parameters, whose numerical values are given in Section 3, have been taken from [17]. Note that the expectation values are scaled by factor of $2m_B$ over those given in some literature due to our different normalization of the meson wavefunctions. It is trivial to check that both conventions yield the same values for physical observables.

We wish to draw the reader's attention to the fact that with changing x , the interference pattern between the SM box and the SUSY box changes. For example, if only $(\delta_{LL}^{13})^2$ is nonzero, there is a constructive interference with the SM box if $x < 2.1$, and a destructive interference otherwise. This is just because c_1 changes sign as we go to higher values of x . Thus, if one chooses $(\delta_{LL}^{13})^2$ to be real, β_{eff} goes down from β for low x , and goes up for high x . Near the crossover region, the SUSY contribution almost vanishes, so one can in principle have large δ parameters. We do not analyze these regions since they smell of a fine-tuning, but one should keep this point in mind.

2.3 R-Parity violating SUSY

R-parity is a global quantum number, defined as $(-1)^{3B+L+2S}$, which is +1 for all particles and -1 for all superparticles. In the minimal version of supersymmetry and some of its variants, R-parity is assumed to be conserved *ad hoc*, which prevents single creation or annihilation of superparticles. However, models with broken R-parity can be constructed naturally, and such models have a number of interesting phenomenological consequences [18, 19]. Some of these R-parity violating models can be motivated from an underlying GUT framework [20].

It is well known that in order to avoid rapid proton decay one cannot have both lepton number and baryon number violating RPV model, and we shall work with a lepton number violating one.

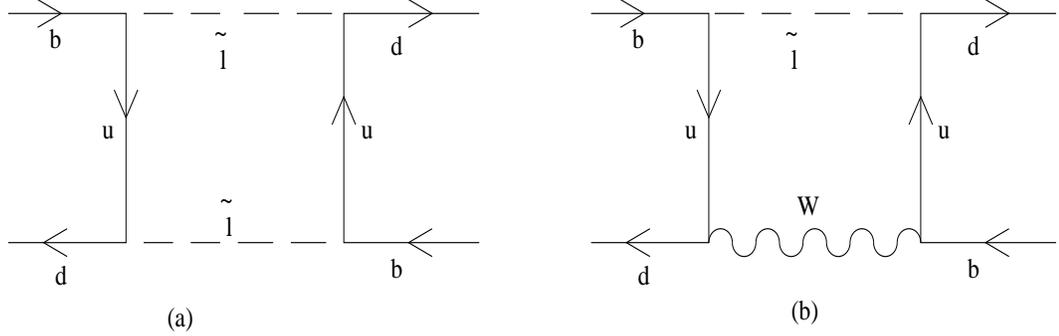


Figure 2: *R*-parity violating contributions to $B^0 - \overline{B^0}$ mixing. Figure (a) corresponds to $L4$, while figure (b) to $L2$ amplitudes (see text for their meanings). One must add the crossed diagrams, as well as the diagrams in (b) where the W is replaced by the charged Higgs or the charged Goldstone. The internal slepton can be of any generation, and so can be the internal charge $+2/3$ quarks, generically depicted as u .

This leads to slepton/sneutrino mediated B decays, and new amplitudes for $B^0 - \overline{B^0}$ mixing with charged sleptons and up-type quarks (and maybe W , charged Higgs and charged Goldstone bosons) flowing inside the loop (see Fig. 2). Since the current lower bound on the slepton mass is weaker than that on squark mass, larger effects within the reach of current round of experiments are more probable in this scenario. We start with the superpotential

$$\mathcal{W}_{\lambda'} = \lambda'_{ijk} L_i Q_j D_k^c, \quad (15)$$

where $i, j, k = 1, 2, 3$ are quark and lepton generation indices; L and Q are the $SU(2)$ -doublet lepton and quark superfields and D^c is the $SU(2)$ -singlet down-type quark superfield respectively. Written in terms of component fields, this superpotential generates six terms, plus their hermitian conjugates, but for our present purpose the only relevant term is

$$\mathcal{L}_{\mathcal{R}} \supset -\lambda'_{ijk} \tilde{e}_L^i \tilde{d}_R^k u_L^j + H.c. \quad (16)$$

With such a term, one can have two different kind of boxes, shown in Fig. 2, that contribute to $B^0 - \overline{B^0}$ mixing: first, the one where one has two sleptons flowing inside the loop, along with two up-type quarks [21], and secondly, the one where one slepton, one W (or charged Higgs boson or Goldstone boson) and two up-type quarks complete the loop [9]. It is obvious that the first amplitude is proportional to the product of four λ' type couplings, and the second to the product of two λ' type couplings times G_F . We call them $L4$ and $L2$ boxes, respectively, for brevity.

The effective Hamiltonian for the $L4$ boxes reads (no sum over k)

$$\mathcal{H}_{L4} = \frac{(\lambda'_{ik1} \lambda'_{ik3})^2}{32\pi^2 \tilde{m}_l^2} I \left(\frac{m_{qk}^2}{\tilde{m}_l^2} \right) \tilde{O}_1 \quad (17)$$

where \tilde{m}_l is the slepton mass, and \tilde{O}_1 has been defined in eq. (9). Actually, there can be different up-type quarks and different generations of sleptons flowing in the box, but that makes the

Hamiltonian proportional to the product of four λ' couplings, which we avoid for simplicity. The function $I(x)$ is given by

$$I(x) = \frac{1 - x^2 + 2x \log x}{(1 - x)^3}. \quad (18)$$

For $L2$ boxes, there are three different types of amplitudes in the Feynman gauge: involving, along with a slepton, a W , a charged Higgs boson, or a charged Goldstone boson. The sum is given by [9]

$$\mathcal{H}_{L2} = \frac{G_F \lambda'_{ik1} \lambda'_{ik3}}{4\sqrt{2}\pi^2} V_{u_k b}^* V_{u_k d} \left[(1 + \cot^2 \beta) x_k^2 J(x_k) + I(x_k) \right] O_4 \quad (19)$$

where $x_k = m_{u_k}^2 / m_{\tilde{m}_l}^2$,

$$J(x) = \frac{-2(x - 1) + (x + 1) \log x}{(x - 1)^3} \quad (20)$$

and $\cot \beta = v_d / v_u$, the ratio of the vacuum expectation values of the two Higgses that give mass to the down- and the up-type quarks respectively (not to be confused with the phase of V_{td}). The Hamiltonian is slightly modified if one has different up-type quarks u_k and u_l in the box:

$$\mathcal{H}_{L2} = \frac{G_F \lambda'_{ik1} \lambda'_{ip3}}{4\sqrt{2}\pi^2} V_{u_k b}^* V_{u_p d} \left[(1 + \cot^2 \beta) x_k x_p K(x_k) + L(x_k) \right] O_4 \quad (21)$$

where we have assumed $m_k > m_p$; if not, the arguments of K and L are to be replaced by x_p . The functions are given by

$$\begin{aligned} K(x) &= \frac{x - 1 - \log x}{(x - 1)^2}, \\ L(x) &= 1 - xK(x). \end{aligned} \quad (22)$$

Note that one can have an imaginary part in the amplitude when the internal quarks are both light, but we neglect that effect for our present purpose³. Also note that in the light of LEP data which definitely favours $\tan \beta \geq 2 - 3$ [23], the Goldstone contributions are dominant over the charged Higgs contributions, which are suppressed by $\cot^2 \beta$. In deriving the above expressions, we have assumed all scalars flowing inside the box to have equal mass.

In general both the λ' type couplings can have a phase, but one of them can be absorbed in the definition of the slepton propagator, so it is enough to consider the one remaining phase.

It is easy to see that the relevant equations (4)-(6) need to be slightly modified to include two NP amplitudes when same quarks flow in the loop; this being a trivial exercise, we do not show the formulae explicitly. For small values of the corresponding λ' couplings, H_{L2} dominates H_{L4} , but the role may get reversed for large values of the product coupling. Thus, one gets two bands in the RPV coupling versus ΔM_d plane, and our bounds correspond to the outer band. There is no such complication when only the $L2$ amplitude is present.

3 Numerical Inputs

The important numerical inputs used in our work is shown in Table 1. A few points are to be noted.

³Even from dimensional arguments, the ratio of the imaginary and the real parts of the mixing amplitude should at most be of the order of m_b^2 / m_t^2 [22], and hence the phase introduced by such an imaginary part can be neglected in the analysis.

Quantity	Value	Remarks
m_B	5.2794 GeV	[24]
m_W	80.423 GeV	[24]
ΔM_d	$0.502 \pm 0.007 \text{ ps}^{-1}$	[25]
$\sin(2\beta_{eff})$	0.681-0.784	[26], at 1σ C.L.
β_{eff}	21.4° - 26.0°	Assumed to be $< 90^\circ$
γ	44° - 72°	[27], at 1σ C.L.
$m_t^{MS}(m_t^{MS})$	166 GeV	[27]
$m_b^{MS}(m_b^{MS})$	4.23 GeV	[6]
$m_b(m_b)$	4.6 GeV	
$m_d(m_b)$	5.4 MeV	
η_B	0.55 ± 0.01	[27]
$f_B\sqrt{B_B}$	$230 \pm 28 \pm 28 \text{ MeV}$	[27]
$\alpha_s(m_Z)$	0.1172 ± 0.002	[24]
$ V_{td} \times 10^3$	6.3 - 9.6	[28], at 95% C.L.
$ V_{ub} \times 10^3$	2.49 - 4.55	[28], at 95% C.L.

Table 1: Input parameters used for the numerical analysis.

Unless shown in the table, we have not taken the experimental uncertainty of a quantity into account. For example, we have used the central values of the CKM elements, except that of V_{td} and V_{ub} , in our analysis. They are taken from Ref. [28], extracted without the world average of $\sin(2\beta)$. This is justified since now $\sin(2\beta)$ itself has a NP contribution, and so should not be used to extract the values of the CKM parameters.

V_{td} is determined from ΔM_d ; thus, the SM fit for V_{td} no longer works when there is one or more NP amplitudes to the box. Since the CP asymmetries are controlled by phases not all of which are in the CKM matrix, the usual argument of the so-called Universal Unitarity Triangle (UUT) [29] does not hold. In essence V_{td} becomes a free parameter, only controlled by the unitarity of the CKM matrix. To take into account this feature, we have taken the 95% confidence limit (CL) for V_{td} , extracted without the global average of $\sin(2\beta)$, for our analysis [28]. The same holds for V_{ub} , which contains the phase γ . As pointed out by [6], γ also becomes a free parameter. We address this issue by keeping the range of γ within 1σ CL quoted by [27], while V_{ub} is varied over its 95% CL range. This, we have checked, essentially covers the whole region generated by a narrower range of V_{ub} and a wider range of γ covering 0 to 2π , with the constraint that the three-generation CKM matrix is unitary (*i.e.*, the unitarity triangle should close). The bounds are of the same order, but slightly more conservative for the former case.

The imaginary part of δs crucially depends on the choice of β . For our analysis, we have varied $\sin(2\beta)$ between 0 and 1, putting a special emphasis on those values which make $\beta \approx \beta_{eff}$.

The most important parameter for the RPC SUSY analysis is the average squark mass, which we fix at 500 GeV, effectively neglecting the splitting caused by the SU(2) D-term. The gluino mass is varied between $0.3 < m_{\tilde{g}}^2/m_{\tilde{q}}^2 < 4.0$. The bounds more or less scale with the squark mass, as can be seen from eqs. (8) and (10). This, however, leaves out some models, *e.g.*, those with an extremely light gluino.

For RPV SUSY analysis, we take all sleptons to be degenerate at 100 GeV. The bounds on the product couplings scale as the square of the slepton mass. The charged Higgs is also assumed to be at 100 GeV. As is discussed below, the precise value of the Higgs mass is not a crucial input.

The value of $\tan\beta$ (the SUSY parameter) is fixed at 3, compatible with the lower bound of the recent LEP analysis [23]. A glance at eqs. (19) and (21) should convince the reader that the bounds are not very sensitive to the exact value of $\tan\beta$, unless it is small, since the Goldstone contributions independent of β control the show. We have explicitly checked the robustness of the bounds with the variation of $\tan\beta$.

The B -parameters have been taken from [17], and at $\mu = m_b$, read

$$B_1 = 0.87(4)^{+5}_{-4}, B_2 = 0.82(3)(4), B_3 = 1.02(6)(9), B_4 = 1.16(3)^{+5}_{-7}, B_5 = 1.91(4)^{+22}_{-7}. \quad (23)$$

The low-energy Wilson coefficients can be found in [6].

4 Results for RPC SUSY

$\sin 2\beta$	x	$\sqrt{ (\delta_{LL}^{13})^2 }$	$\sqrt{ (\delta_{LR}^{13})^2 }$	$\sqrt{ \delta_{LL}^{13} \delta_{RR}^{13} }$	$\sqrt{ \delta_{LR}^{13} \delta_{RL}^{13} }$
1.0	0.3	0.046	0.021	—	—
	1.0	0.099	0.023	—	—
	2.0	0.27	0.026	—	—
	4.0	—	0.031	—	—
0.732	0.3	0.017	0.0075	0.0030	0.0040
	1.0	0.036	0.0080	0.0033	0.0068
	2.0	0.10	0.0095	0.0039	0.0099
	4.0	0.090	0.012	0.0048	0.015
0.5	0.3	—	—	0.0078	0.011
	1.0	—	—	0.0088	0.018
	2.0	—	—	0.010	0.026
	4.0	0.235	—	0.012	0.039

Table 2: Bounds on the δ^{13} parameters when they are all real.

The real and imaginary parts of the SUSY amplitude being established through the real and imaginary parts of the corresponding δ s, it is easy to find the limits on those real and imaginary parts. We perform a scan on the complete range of V_{td} and $f_B\sqrt{B_B}$, as well as on the SUSY phase ϕ , where generically $\delta^2 = |\delta^2| \exp(-2i\phi)$, over the range 0 to π . We demand that ΔM_d and $\sin(2\beta_{eff})$ should lie between the values specified in Table 1. The results are obtained for various values of β , even for extreme values like $\sin(2\beta) = 0$ or 1.

Tables 2 and 3 summarize our results. In table 2, we show the bounds on various δ parameters (and their combinations) when they are real, assuming only one to be nonzero at a time. Note how the interference pattern changes for δ_{LL}^{13} with x ; it is constructive for small x , and destructive for large x . Near $x = 2$, the crossover occurs, so one can have a large value for δ_{LL}^{13} . A look at the respective Wilson coefficients and hence the interference pattern should help the reader to

$\sin 2\beta$	x	$\sqrt{ Re(\delta_{LL}^{13})^2 }$	$\sqrt{ Im(\delta_{LL}^{13})^2 }$	$\sqrt{ Re(\delta_{LR}^{13})^2 }$	$\sqrt{ Im(\delta_{LR}^{13})^2 }$
1.0	0.3	0.046	0.079	0.021	0.035
	1.0	0.10	0.18	0.022	0.039
	2.0	0.27	0.27	0.026	0.046
	4.0	0.23	0.30	0.031	0.061
0.732	0.3	0.077	0.078	0.028	0.030
	1.0	0.16	0.16	0.035	0.038
	2.0	0.27	0.28	0.038	0.040
	4.0	0.24	0.25	0.051	0.054
0.5	0.3	0.081	0.070	0.031	0.027
	1.0	0.17	0.15	0.038	0.032
	2.0	0.31	0.23	0.046	0.038
	4.0	0.28	0.19	0.053	0.044
0.0	0.3	0.084	0.048	0.034	0.021
	1.0	0.18	0.10	0.040	0.026
	2.0	0.27	0.28	0.054	0.035
	4.0	0.30	0.24	0.060	0.032

Table 3: Bounds on the real and the imaginary parts of δ_{LL}^{13} and δ_{LR}^{13} .

understand the missing entries. Without the CP-asymmetry constraint, nontrivial entries would occur everywhere [5].

Table 3 shows the real and imaginary parts for δ_{LL}^{13} and δ_{LR}^{13} for several values of $\sin(2\beta)$. Due to the presence of the nonzero SUSY phase, all entries are nonvanishing. Also, the rise of δ_{LL}^{13} near $x = 2$ is less pronounced. We do not show the entries for negative values of $\sin(2\beta)$; they are more or less equal with their counterparts for positive $\sin(2\beta)$. There are a few exceptions, which do not change the general result.

We caution the reader that these bounds are the most conservative ones at the particular benchmark points that we have chosen, but by no means signify an impossibility of having larger δ s at other points in the SUSY parameter space. Particularly, note that these bounds scale with the squark mass.

5 Results for RPV SUSY

We explicitly assume that the contributions coming from the RPC SUSY sector vanish if there is nonzero RPV interactions. This will be justified *post hoc* when we discuss a sample case where both are present, and there is a possibility of cancellation.

The strategy is the same as that adopted for the previous subsection. The phase of the RPV product coupling is varied between 0 and 2π while the magnitude of the coupling is assumed to be only positive. The range of the scan is kept between the direct product limits, *i.e.*, the limit one obtains when one multiplies the individual limits for the two λ' components. This limit, as can be easily checked [30], is most lenient for the third slepton generation. Only two of the relevant product couplings have been bounded from other sources: the product $\lambda'_{i13}\lambda'_{i31}$ has a very stringent

$\lambda'\lambda'$ combination	$Re(\lambda'\lambda')$	$Im(\lambda'\lambda')$	Direct bound
(i31)(i33)	1.8×10^{-3}	1.7×10^{-3}	0.202
(i21)(i23)	1.2×10^{-3} (0.022)	$1.2(1.4) \times 10^{-3}$	0.27
(i11)(i13)	$2.4(2.6) \times 10^{-3}$	$2.5(2.6) \times 10^{-3}$	2.7×10^{-3} (*)
(i11)(i23)	0.016	0.016	0.057
(i11)(i33)	0.026 [†]	0.026 [†]	0.026
(i21)(i13)	$2.5(3.0) \times 10^{-4}$	2.8×10^{-4}	0.057
(i21)(i33)	0.098	0.1	0.23
(i31)(i23)	1.4×10^{-4}	$1.35(1.6) \times 10^{-4}$	0.26

Table 4: Upper limits on the real and the imaginary parts of the relevant $\lambda'\lambda'$ couplings for $\sin(2\beta) = 0.732$. The numbers in parenthesis show the maximum possible value, for some other $\sin(2\beta)$. All the products, except the one marked with a dagger ([†]), show improvements over the corresponding bounds obtained from direct product of the bounds of the relevant λ' s (for $i = 3$), shown in the fourth column. All the numbers in the fourth column are from [30], except the one marked with asterisk, which is from [12]. The bounds on the imaginary parts are new.

bound from tree-level $B^0 - \overline{B^0}$ mixing, of the order of 10^{-8} (and hence we do not discuss this product further), and the product $\lambda'_{i11}\lambda'_{i13}$ has been constrained from the measured branching ratio and CP asymmetries in the $B \rightarrow \pi^+\pi^-$ channel [12] (marked with an asterisk in Table 4). There is no bound on the imaginary parts of the couplings.

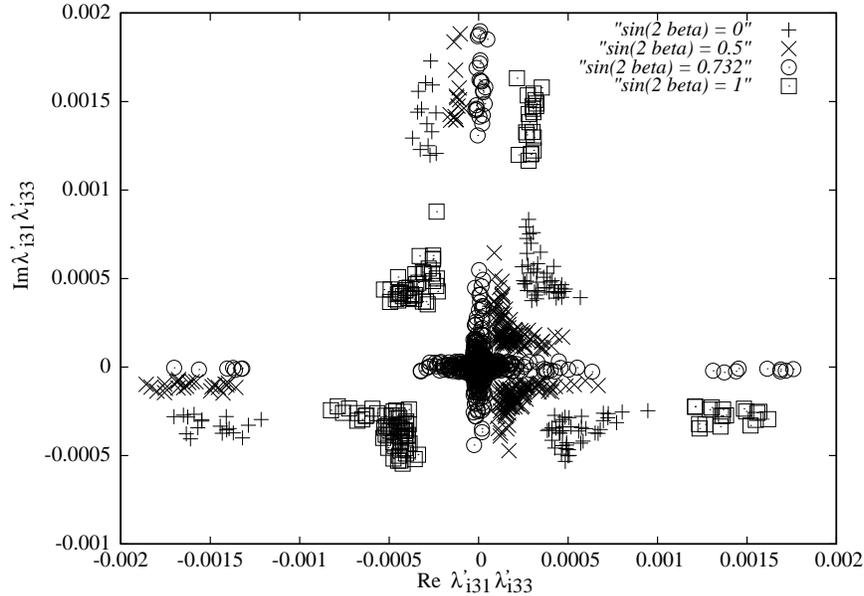


Figure 3: The real and imaginary parts of $\lambda'_{i31}\lambda'_{i33}$ for various values of $\sin(2\beta)$.

Table 4 summarizes our results. Note that almost all the products, apart from those two discussed above, and the one marked with a dagger, have been improved, some by orders of magnitude. Six among these eight entries have been considered by the authors of [9]. One may note that they obtained bounds which have the same orders of magnitude to the ones that we get. However, there are several ways in which we have improved upon their calculation, apart from

taking the updated data as input. These improvements are: (i) imposition of the CP-asymmetry constraint, which was not available at their time (this helps us to obtain the bounds on the imaginary parts of the couplings); (ii) incorporation of the NLO QCD corrections, which are, anyway, expected to be small — at least they should not change the numbers by, say, a factor of two; (iii) scan over the full range of the input parameters, as we have already discussed; and (iv) consideration of the SM contribution, and the possibilities of interference between SM, $L2$ and $L4$ amplitudes. Also note that they have computed the bounds for $\tan\beta = 1$.

Our bounds are more or less insensitive to the precise value of $\sin(2\beta)$ (the UT angle); the exceptions are shown in the table. Another feature is that real and imaginary parts have almost the same bounds. Figure 3 highlights this nature in detail for the product coupling $\lambda'_{i31}\lambda'_{i33}$.

Let us try to understand this figure. There are three main regions, the first one (like a smeared cross) encompassing the origin, the second one, divided into four almost symmetrical fragments, around $|Re(\lambda'_{i31}\lambda'_{i33})| = 5 \times 10^{-4}$, and a third fragmented region about $|\lambda'_{i31}\lambda'_{i33}| = 15 \times 10^{-4}$. It is the outermost third region which gives the bound. It is easy to explain the origin of these three regions, and sheds light on the role of $\sin(2\beta)$ in determining the parameter space.

The first one is governed by the SM, and the difference between the experiment and the theory is filled up by RPV. Note that this region has points only for $\sin(2\beta) = 0.732$; thus, SM is allowed without any NP. There is a satellite region, for $\sin(2\beta) = 0.5$, where SM is not allowed (from the A_{CP} constraint), and RPV fills in to generate the necessary CP asymmetry. The second region includes points for extreme values of $\sin(2\beta)$: 0 or 1. SM is not allowed, and one needs a greater role from RPV to obtain the observed CP asymmetry. Though both $L2$ and $L4$ boxes are allowed for this RPV coupling, the $L2$ box dominates here. However, they come with opposite sign, so there is a region where they interfere destructively (particularly with the increase of the coupling), and RPV contribution may even go to zero, leaving only SM. This generates the third region and explains why one has points even for $\sin(2\beta) = 0.732$.

It is a possibility that both RPC and RPV SUSY are present, however pathological that may seem. Even in this extreme case the bounds are never changed by orders of magnitude, unless there is a very precise fine-tuning. For example, with only δ_{LL}^{13} and $\lambda'_{i21}\lambda'_{i33}$ present (this particular combination is chosen since both of them have comparable upper bounds), the bounds are relaxed to $\sqrt{|Re(\delta_{LL}^{13})^2|} = 0.22$, $\sqrt{|Im(\delta_{LL}^{13})^2|} = 0.26$, $Re(\lambda'_{i21}\lambda'_{i33}) = 0.12$, $Im(\lambda'_{i21}\lambda'_{i33}) = 0.11$, for $x = 1$ and $\sin(2\beta) = 0.732$. Thus we have reasons to be confident about these bounds. This is more so if the limits have different orders of magnitude.

$\lambda'\lambda'$ combination	Decay channels	$\lambda'\lambda'$ combination	Decay channels
(i31)(i33)	$b \rightarrow d\bar{\ell}_i\ell_i, b \rightarrow d\bar{\nu}_i\nu_i$	(i21)(i23)	$b \rightarrow c\bar{c}d, b \rightarrow s\bar{s}d$
(i11)(i23)	$b \rightarrow c\bar{u}d, b \rightarrow s\bar{d}d$		$b \rightarrow d\bar{\ell}_i\ell_i, b \rightarrow d\bar{\nu}_i\nu_i$
(i21)(i13)	$b \rightarrow u\bar{c}d, b \rightarrow d\bar{s}d$	(i11)(i13)	$b \rightarrow u\bar{u}d, b \rightarrow d\bar{d}d$
			$b \rightarrow d\bar{\ell}_i\ell_i, b \rightarrow d\bar{\nu}_i\nu_i$

Table 5: The possible decay channels of the B meson driven by the different RPV product couplings. The final state mesons are not shown explicitly. For the semileptonic decays, the outgoing leptons must be of the same generation, denoted by i .

Most of these product couplings contribute to various B-decay channels, both nonleptonic and semileptonic. They are enlisted in Table 5. It is easy to check which mesons and leptons come out

in the final stage. The present e^+e^- B factories, as well as the future hadronic machines, should put tighter constraints on these RPV product couplings.

6 Summary and Conclusions

In this paper we have enlisted the constraints on the real and the imaginary parts of the FCNC parameters of both R-parity conserving and R-parity violating SUSY, coming from $B^0 - \overline{B}^0$ mixing. For RPC SUSY, these are the conventional δ_{13}^d parameters of different chiralities. The same analysis was performed by [6]; our results differ slightly from theirs due to two reasons. We have performed a scan over all SM quantities, including V_{td} and $f_B^2 B_B$, and our range of scan for γ is different from theirs. Weaker constraints on these δ parameters can also be derived from the radiative decay $b \rightarrow d\gamma$.

For the RPV SUSY scenario, the FCNC parameters are the λ' type lepton-number violating couplings. One needs a product of two such couplings to generate $B^0 - \overline{B}^0$ oscillation. There can be several such products, depending on the choice of the quarks and sleptons flowing inside the box, whose real and imaginary parts have been bounded from the experimental data. The bounds on the real parts update the work of [9] while the bounds on the imaginary parts are derived for the first time in this paper.

If we ever find a signal for NP in B factories, how can we be certain that it indeed comes from SUSY? There are three steps to ascertain that. First, sort out those channels which show an abnormality. Second, Try to find the model which can explain these anomalies. And third, check whether there are other channels where one may expect to see an anomaly, and whether the anomaly may be present in the data. If there is a prospective channel, one should look for it. Confirmation of the nature of the NP is never possible without the study of such correlated signals. Such correlated signals may even be the direct production of new particles, *e.g.*, in the Large Hadron Collider. A possible discriminating signal between RPC and RPV SUSY is the fact that the RPV version may be highly flavor-specific, and so one would expect the absence of anomaly in such channels which may be affected in a more flavor-blind model such as the RPC SUSY.

For the help of our experimentalist colleagues, we enlist the possible decay modes of the B meson which are driven by the RPV product couplings discussed here. A careful study of the possible channels in present and future B factories should be able to put tighter constraints on the parameter space. Particularly, the proposed higher luminosity e^+e^- B factories should make both $\sin(2\beta)$ and ΔM_d precision observables, and the bounds are expected to be improved by at least one order of magnitude, if we do not discover SUSY by that time.

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