Physics Letters B 706 (2012) 379–383

Contents lists available at SciVerse ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

$B \to \tau\nu$: Opening up the charged Higgs parameter space with R-parity violation

Roshni Bose, Anirban Kundu *

Department of Physics, University of Calcutta, 92, Acharya Prafulla Chandra Road, Kolkata 700 009, India

ARTICLE INFO

Article history:
Received 11 October 2011
Received in revised form 9 November 2011
Accepted 20 November 2011
Available online 23 November 2011
Editor: T. Yanagida

ABSTRACT

The theoretically clean channel $B^+ \to \tau^+\nu$ shows a close to 3σ discrepancy between the Standard Model prediction and the data. This in turn puts a strong constraint on the parameter space of a two-Higgs doublet model, including R-parity conserving supersymmetry. The constraint is so strong that it almost smells of fine-tuning. We show how the parameter space opens up with the introduction of suitable R-parity violating interactions, and release the tension between data and theory.

© 2011 Elsevier B.V. Open access under CC BY license.

1. Introduction

The purely leptonic decay $B^+ \to \tau^+\nu$ has generated a lot of interest in recent times because both BaBar and Belle Collaborations found a sizable discrepancy with the Standard Model (SM) prediction, which is quite clean and robust. The world average is [1]

$$Br(B \to \tau\nu) = (16.8 \pm 3.1) \times 10^{-5}, \tag{1}$$

while the theoretical prediction is

$$Br(B \to \tau\nu)_{SM} = \left(7.57^{+0.98}_{-0.61}\right) \times 10^{-5}, \tag{2}$$

which gives a tension at the level of 2.8σ [1]. The ratio of experimental and SM-expected branching fraction is approximately $2.22^{+0.50}_{-0.45}$. The theoretical uncertainty comes from the B meson decay constant f_B and the Cabibbo–Kobayashi–Maskawa (CKM) matrix element V_{ub}.

Using the Lattice QCD result of [2]

$$f_{B_s} = (231 \pm 3 \pm 15)\ \text{MeV},$$
$$f_{B_s}/f_{B_d} = 1.235 \pm 0.008 \pm 0.033, \tag{3}$$

an SM-only explanation would require [1]

$$|V_{ub}| = (5.10 \pm 0.59) \times 10^{-3}, \tag{4}$$

which is clearly inconsistent with the indirect determination of V_{ub} from the sides of the Unitarity Triangle (UT) [1],

$$|V_{ub}|_{\text{indirect}} = (3.49 \pm 0.13) \times 10^{-3}, \tag{5}$$

or the average of direct inclusive ($B \to X_u\ell\nu$) and exclusive ($B \to \pi\ell\nu$) measurements [3],

$$|V_{ub}|_{\text{measured}} = (3.92 \pm 0.09 \pm 0.45) \times 10^{-3}. \tag{6}$$

This gives rise to the speculation that beyond-SM (BSM) physics might be at work here. One of the possibilities is a possible new physics in B^0–$\bar{B}^0$ mixing so that the indirect measurement of V_{ub} is not its SM value. Apart from that, the first candidate for such BSM physics is the charged Higgs boson H^+ of the two-Higgs doublet models (2HDM) [4], or of any supersymmetric model. The charged Higgs contribution has a destructive interference with the SM W-mediated contribution, so the solution comes only as a narrow band, centred roughly about

$$m_H\ (\text{GeV}) \approx 3.3 \tan\beta, \tag{7}$$

where $\tan\beta$ is the usual ratio of the two vacuum expectation values [5]. It has been argued that the solution looks rather unnatural and almost smells of fine-tuning [6]. This solution, for the 2HDM model type II, also suffers serious tension from processes like $B \to D\ell\nu$, the ratio $Br(K \to \mu\nu)/Br(\pi \to \mu\nu)$, $b \to s\gamma$, $Z \to b\bar{b}$, and the neutral B meson mass differences ΔM_d and ΔM_s. As was shown in [6], the fine-tuned region disappears when one takes all B-physics data into account at 95% confidence limit (CL). However, it was recently shown in [7] that a Minimal Flavour Violating 2HDM has a better agreement to these observables.

Models which embed the 2HDM, like supersymmetry (SUSY), have also been studied. The conclusions, however, are not very enthusiastic [5]. The reason is that in R-parity conserving SUSY (the definition of R-parity is given later), which is phenomenologically attractive because of its cold dark matter candidate, the SUSY effects to $B \to \tau\nu$ appear only as one-loop diagrams with heavy electroweak gauginos and sleptons in the loop. Thus, the new amplitudes open up the parameter space only marginally [5].

* Corresponding author.
E-mail addresses: 123.roshni@gmail.com (R. Bose), akphy@caluniv.ac.in (A. Kundu).

0370-2693 © 2011 Elsevier B.V. Open access under CC BY license.
doi:10.1016/j.physletb.2011.11.044

R-parity violating (RPV) SUSY has also come up as another interesting option [8–10]. The lepton and baryon numbers, L and B respectively, are good symmetries of the Standard Model but *ad hoc* symmetries of the minimal supersymmetric SM, in the sense that one can write L and B violating terms in the superpotential. However, conservation of both L and B leads to a Z_2 symmetry called R-parity, and defined as

$$R_p = (-1)^{3B+L+2S}, \tag{8}$$

where S is the spin of the particle. By definition, $R_p = +1$ for particles and $R_p = -1$ for superparticles, and we demand R_p to be a good symmetry of the superpotential so that the Z_2 symmetry leads to a dark matter candidate. On the other hand, if R_p is not a good symmetry, the signatures change drastically, because all superparticles, including the LSP, can decay inside the detector.

There can be 45 R_p-violating (RPV) couplings in the superpotential coming from the renormalizable terms

$$W_{R_{\not p}} = \epsilon_{ab}\left(\frac{1}{2}\lambda_{ijk}L_i^a L_j^b \bar{E}_k + \lambda'_{ijk}L_i^a Q_j^b \bar{D}_k\right) + \frac{1}{2}\epsilon_{pqr}\lambda''_{ijk}\bar{U}_i^p \bar{D}_j^q \bar{D}_k^r, \tag{9}$$

where L, Q, E, U and D stand for lepton doublet, quark doublet, lepton singlet, up-type quark singlet, and down-type quark singlet superfields respectively; i, j, k are generation indices that can run from 1 to 3; $a, b = 1, 2$ are SU(2) indices, and $p, q, r = 1, 2, 3$ are SU(3) indices; and λ_{ijk} (λ''_{ijk}) are constructed to be antisymmetric in i and j (j and k). The phenomenology of RPV supersymmetry, including the collider signatures and bounds on these couplings, may be found in [11]. Apart from the trilinear terms, there can be bilinear R-parity violating terms of the form of $-\mu_i L_i H_2$, where H_2 is the superfield that gives mass to charged leptons and down-type quarks, in $W_{R_{\not p}}$. We assume these bilinear terms to be zero at the weak scale. This also relaxes the possible constraints coming from the neutrino masses and mixing angles in presence of the bilinear terms. However, even some trilinear combinations like $\lambda^{(\prime)}_{ikl}\lambda^{(\prime)}_{jlk}$ can generate nonzero entries for the ij-th element of the neutrino mass $\mathcal{M}_\nu$ [11]. Putting the bilinear terms equal to zero means that we choose a particular basis in the $\{H_2, L_i\}$ space. The sneutrino vacuum expectation values need not be zero in this basis, but that is not important for our case.

In this Letter, we will try to use RPV SUSY from a different point of view. We will not constrain the RPV couplings from $B \to \tau\nu$; this has already been done in [9,10], and there are other comparable or stricter bounds [12,13]. We will rather see how much the charged Higgs parameter space in the m_H-$\tan\beta$ plane opens up because of a new tree-level contribution coming from RPV SUSY. Considering the results of [6], such a study is of serious importance. We will also consider the possible effects of the complex phase of the RPV couplings. It will be shown that with some couplings, the parameter space substantially opens up and there is no longer any 'fine-tuning'; with some other couplings, the effect is rather small because they are too tightly constrained. Experimentally, this means that if the $B \to \tau\nu$ data remains anomalous, and we find the charged Higgs at some other point than that allowed by the narrow fine-tuning band, it will indicate another new tree-level contribution; RPV SUSY is a prime candidate for this.

The Letter is arranged as follows. In Section 2, we compile the relevant formulae. The numerical analysis is taken up in Section 3, and we summarize and conclude in Section 4.

2. Relevant expressions

The decay width of $B \to \tau\nu_\tau$, in the SM, is given by

$$\Gamma(B \to \tau\nu_\tau) = \frac{1}{8\pi}G_F^2|V_{ub}|^2 f_B^2 m_\tau^2 m_B\left(1 - \frac{m_\tau^2}{m_B^2}\right)^2, \tag{10}$$

which, in the presence of a charged Higgs, is modified by a multiplicative factor,

$$\Gamma(B \to \tau\nu_\tau)_{2HDM} = \Gamma(B \to \tau\nu_\tau)_{SM} \times \left(1 - \frac{m_B^2}{m_H^2}\tan^2\beta\right)^2. \tag{11}$$

If R-parity is violated (there are two leptons in the final state, so we will consider only L-violating interactions, *i.e.*, all λ'' couplings are assumed to be zero), there are new contributions to the amplitude. These contributions, as has been pointed out in [9], can either be squark mediated (with a generic form of $\lambda'\lambda'$), or slepton mediated (with a generic form of $\lambda\lambda'$). Each individual coupling can in general be complex. While it is possible to absorb the phase of one coupling by redefining the phase of the propagating sfermion, the phase of the second coupling cannot be absorbed, and so in general this quadratic product of the RPV-couplings is complex.

The ATLAS and CMS experiments at the Large Hadron Collider have already ruled out squarks below 800 GeV. Assuming a universal scalar mass M_0, one can safely integrate out the propagating squark or slepton fields, and arrive at a four-fermion interaction analogous to the Fermi interaction. As is the standard practice, we will assume a hierarchical structure of the RPV couplings, so effectively only one product is nonzero at the weak scale. This might not be the case if the RPV couplings are defined in the weak basis and one rotates them by a CKM-like mechanism to get the relevant couplings in the mass basis [14], but the qualitative results do not change much. Even in the mass basis, the RPV couplings λ_{ijj} ($i \neq j$) or λ'_{ijj} are severely restricted, as they are possible sources of neutrino Majorana mass terms.

One also notes that with RPV, the final state neutrino can have any flavour, depending on the couplings. If it is a ν_τ, the $R_{\not p}$-amplitude will add coherently with the SM and 2HDM amplitudes; if it is ν_e or ν_μ, the addition will be incoherent. For the latter case, the weak phase of the RPV coupling will not matter.

The relevant four-fermion operators for the decay $B \to \tau\nu$ may be obtained by integrating the sfermion field out in Eq. (9). The expression reads [12,13]

$$\mathcal{H}_{R_{\not p}} = A_{jklm}\big(\bar\nu_j(1+\gamma_5)\ell_m\big)\big(\bar d_l(1-\gamma_5)u_k\big) - \frac{1}{2}B_{jklm}\big(\bar\nu_j\gamma^\mu(1-\gamma_5)\ell_m\big)\big(\bar d_l\gamma_\mu(1-\gamma_5)u_k\big), \tag{12}$$

where

$$A_{jklm} = \sum_{i=1}^{3}\frac{\lambda^*_{ijm}\lambda'_{ikl}}{4M^2_{\tilde\ell_{iL}}}, \qquad B_{jklm} = \sum_{i=1}^{3}\frac{\lambda'_{mki}\lambda'^*_{jli}}{4M^2_{\tilde d_{iR}}}. \tag{13}$$

We take, keeping the recent ATLAS and CMS results in view, all sfermions to be degenerate at 1 TeV. The bounds scale with $M^2_{\tilde f}$, as is evident from Eqs. (12) and (13). For our case, $k = 1$, and $l = m = 3$.

The contributions to the decay width of $B \to \tau\nu$ are given as

$$\mathcal{M}_{SM+2HDM} = \frac{1}{\sqrt{2}}G_F V^*_{ub}\left(1 - \frac{m_B^2}{m_H^2}\tan^2\beta\right),$$

$$\mathcal{M}_{squark} = \frac{1}{2}B_{j133},$$

$$\mathcal{M}_{slepton} = -A_{j133} \times \frac{m_B^2}{m_\tau(m_b + m_u)} \approx -3.7A_{j133}, \tag{14}$$

Table 1
Bounds on RPV couplings at 2σ with all sfermions degenerate at 1 TeV. $\mathcal{M}_{\rm SM} = G_F|V^*_{ub}|/\sqrt{2}$. We have used a perturbative upper bound of 2 on some of the individual $R_{\not p}$ couplings. The entries marked with ‡ are to be added coherently with the SM amplitude, and those marked with † only for $i = 3$. For details, see text.

$R_{\not p}$ coupling	Bound	$\mathcal{M}_{R_{\not p}}/\mathcal{M}_{\rm SM}$	$R_{\not p}$ coupling	Bound	$\mathcal{M}_{R_{\not p}}/\mathcal{M}_{\rm SM}$
$\lvert\lambda'_{313}\lambda'_{i33}\rvert^{\dagger}$	3.8×10^{-4}	1.5×10^{-3}	$\lvert\lambda_{i23}\lambda'_{i13}\rvert$	0.017	0.29
$\lvert\lambda'_{311}\lambda'_{131}\rvert$	0.16	0.62	$\lvert\lambda_{i13}\lambda'_{i13}\rvert$	0.019	0.35
$\lvert\lambda'_{311}\lambda'_{231}\rvert$	0.95	3.68	$\lvert\lambda_{133}\lambda'_{113}\rvert^{\ddagger}$	8.1×10^{-4}	0.023
$\lvert\lambda'_{311}\lambda'_{331}\rvert^{\ddagger}$	0.04	0.16	$\lvert\lambda_{233}\lambda'_{213}\rvert^{\ddagger}$	1.1×10^{-3}	0.031
$\lvert\lambda'_{312}\lambda'_{332}\rvert^{\ddagger}$	0.04	0.16			
$\lvert\lambda'_{312}\lambda'_{1(2)32}\rvert$	1.2	4.64			

where we have used $m_B = 5.27$ GeV, $m_\tau = 1.777$ GeV, and the running mass $m_b \equiv \overline{m_b}(\overline{m_b}) = 4.22$ GeV, corresponding to a pole mass $m_b(\text{pole}) = 4.63$ GeV.

The branching fraction is

$$Br(B \to \tau\nu) = \frac{1}{4\pi} f_B^2 m_\tau^2 m_B \tau_B \left(1 - \frac{m_\tau^2}{m_B^2}\right)^2 \times \left|\mathcal{M}_{\rm SM+2HDM} + \mathcal{M}_{\rm squark/slepton}\right|^2, \tag{15}$$

where $\tau_B = 1.525$ ps is the lifetime of B^+. Eq. (15) depends on the relative phase θ ($0 \leqslant \theta < 2\pi$) of V^*_{ub}, which by convention is denoted by γ or ϕ_3, and the relevant RPV product coupling. While Eq. (15) is true only for $j = 3$, *i.e.*, when the emitted neutrino is ν_τ, for $j = 1$ or 2 one can simply set $\theta = \pi/2$.

Some of the relevant product couplings also contribute to other processes. Most notable among them are the lepton flavour violating (LFV) decays $B_d \to e^-\tau^+$ and $B_d \to \mu^-\tau^+$. They put a rather tight constraint on the corresponding couplings. Other processes include the decay $B \to \pi\nu\bar{\nu}$, from which the constraints are relatively poor, or LFV decays of the top quark. The relevant expressions can be found in [12,13]. One can summarize the bound on the $\lambda\lambda'$ combination coming from $B_d \to \ell\bar{\ell}'$ decay as

$$\left|\lambda\lambda'\right| = 0.017\left(\frac{Br(B \to \ell\bar{\ell}')}{25\times10^{-6}}\right)^{1/2}\left(\frac{\tilde{M}}{1000}\right)^2\left(\frac{m_b}{4.2}\right)\left(\frac{f_B}{0.2}\right)^{-1}, \tag{16}$$

where the sfermion mass $\tilde{M}$, the b-quark running mass m_b, and the decay constant f_B are all measured in GeV.

3. Numerical analysis

The upper bounds on the relevant RPV couplings λ and λ' are taken from [11], with updates from [15]. Some of the product couplings are bounded by LFV decays of the B_d meson, and the numbers are taken from Eq. (16) (also see [12,13]). As all these bounds are at 2σ, we will show our numerical analysis at that confidence level only.

Neutrino masses put the tightest constraint on some of the RPV couplings. Assuming flat ΛCDM cosmology, the total mass of all SM-like ν and $\bar{\nu}$ species is bounded [16] to be $\sum m_\nu < 0.56$ eV. Taking each entry of the neutrino mass matrix to be of the order of 0.28 eV, and all sfermions degenerate at 1 TeV, one gets some typical bounds [11]:

$$\left|\lambda'_{i11}\right| < 5.3\times10^{-1}, \qquad \left|\lambda'_{i33}\right| < 6.4\times10^{-4}, \qquad |\lambda_{i33}| < 2.7\times10^{-3}. \tag{17}$$

There are also single-coupling 2σ bounds [15],

$$|\lambda_{12k}| < 0.3, \quad |\lambda_{13k}| < 0.3, \quad |\lambda_{23k}| < 0.5, \quad \left|\lambda'_{11k}\right| < 0.3, \quad \left|\lambda'_{21k}\right| < 0.4, \quad \left|\lambda'_{31k}\right| < 0.6, \quad \left|\lambda'_{1k1}\right| < 0.3, \tag{18}$$

where $k = 1, 2, 3$. We need some more couplings on which the 2σ bounds at 1 TeV are rather weak [11],

$$\left|\lambda'_{133}\right| < 1.8, \qquad \left|\lambda'_{2k1}\right| < 1.8, \tag{19}$$

and some more on which the limit comes from the expectation that they remain perturbative at the weak scale. Since this is a matter of choice, we impose a flat cutoff at 2:

$$\left|\lambda'_{132}\right|, \left|\lambda'_{23k}\right|, \left|\lambda'_{33k}\right| < 2. \tag{20}$$

The LFV decays of B_d yield the following bounds:

$$\left|\lambda_{123}\lambda'_{113}\right|, \left|\lambda_{233}\lambda'_{313}\right| < 0.017 \quad \left(B_d \to \mu^-\tau^+\right), \qquad \left|\lambda_{123}\lambda'_{213}\right|, \left|\lambda_{133}\lambda'_{313}\right| < 0.019 \quad \left(B_d \to e^-\tau^+\right). \tag{21}$$

The mass difference for the B^0–$\overline{B}^0$ system, ΔM_d, gives [17]

$$\left|\lambda'_{i1j}\lambda'_{i3j}\right| < 0.04. \tag{22}$$

Table 1 summarizes the best bounds at the 2σ level.

The opening up of the charged Higgs parameter space depends on the ratio

$$\mathcal{R} = \frac{\mathcal{M}_{R_{\not p}}}{\mathcal{M}_{\rm SM}}, \tag{23}$$

as displayed in Table 1. For example, for $\mathcal{R} = 1.5\times10^{-3}$, the change is imperceptible. For other typical values, we refer the reader to Fig. 1.

The figure shows the charged Higgs parameter space for $m_H \in [100:1000]$ GeV and $\tan\beta \in [3:100]$. For the coupling $\lambda'_{311}\lambda'^{*}_{331}$, the amplitude addition is coherent. There is a marginal enhancement on both sides of the pure-2HDM band (only for those RPV amplitudes that add coherently with the SM one, the enhancement is above the 2HDM band; for incoherent additions, only the lower portion of the band might get allowed), and another region with low $\tan\beta$ opens up. This is the region where charged Higgs contribution is insufficient to make up the deficit, but that role is taken up by RPV SUSY. The lower left plot is for $\mathcal{R} = 0.62$, where the relevant coupling is $|\lambda'_{311}\lambda'_{131}|$. The emitted neutrino is ν_e and so the amplitudes add incoherently; the parameter space opens up only on the lower side of the upper edge of pure-2HDM region. Note that the gap between the two allowed regions shrink. For $\mathcal{R} = 4.64$, where the coupling is either $|\lambda'_{312}\lambda'_{132}|$ or $|\lambda'_{312}\lambda'_{232}|$, these two regions merge. The addition being incoherent, the allowed region is again bounded by the upper edge of the 2HDM band.

We do not show the plots for every possible coupling, because the trend is obvious. The two regions merge at about $\mathcal{R} = 1.18$, whether the sum is coherent or incoherent. (However, no RPV couplings with such large $\mathcal{R}$ that can add coherently with the SM amplitudes are allowed.) Thus, the charged Higgs parameter

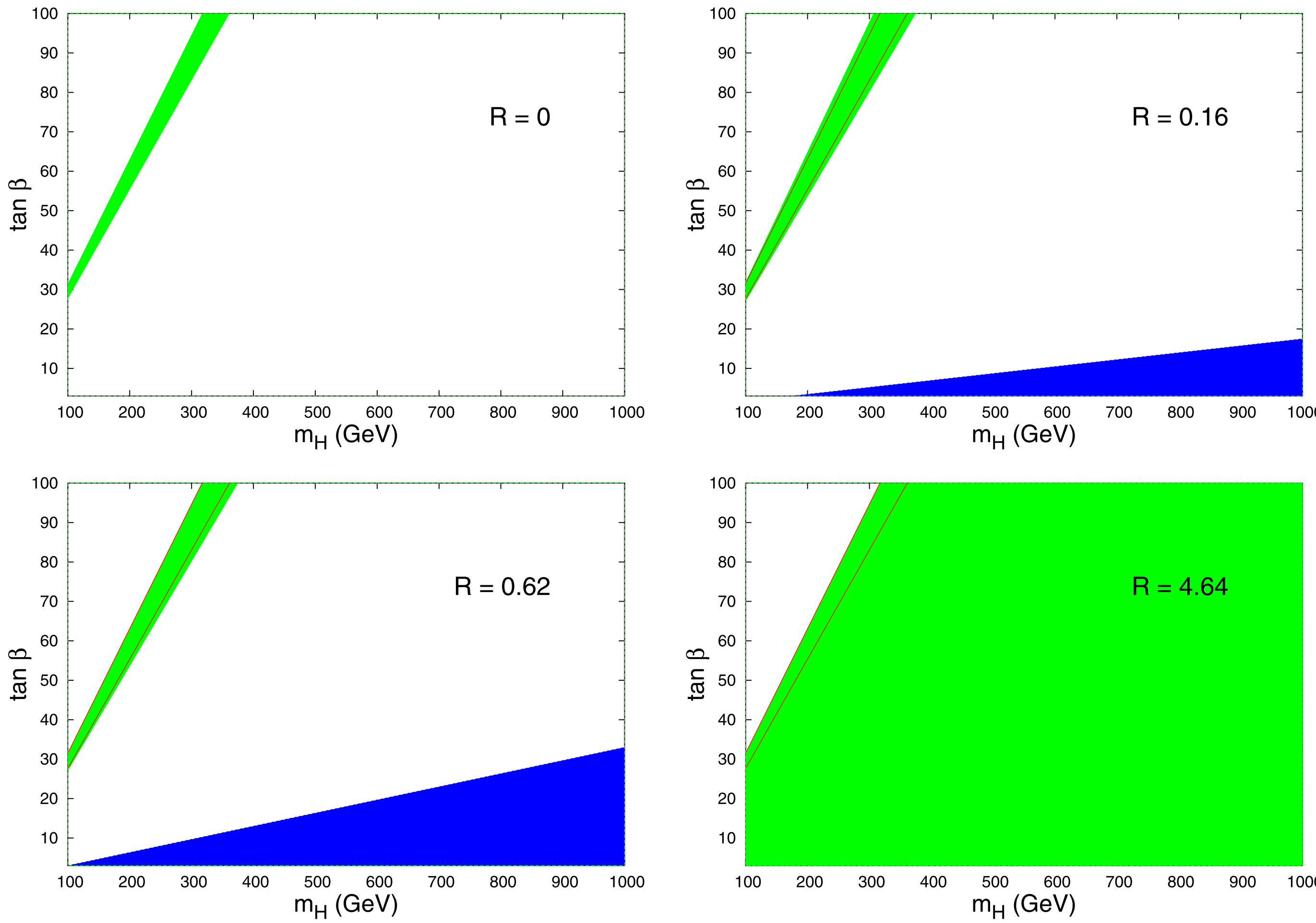

Fig. 1. The allowed parameter space, at 95% CL, for the charged Higgs, as a function of $\mathcal{R}$ defined in Eq. (23). Left upper plot shows the 'only-2HDM' fine-tuned parameter space. For the rest three plots, this region is that between the solid lines. Right upper and left lower plots show the allowed parameter spaces (blue/dark grey and green/light grey shaded regions) for different values of $\mathcal{R}$. The lower right plot is for a large value of the RPV couplings, where the blue/dark grey and green/light grey regions merge.

space opens up significantly, unless the corresponding RPV coupling is very tightly constrained. In particular, the low $\tan\beta$ region becomes allowed and alleviates the tension with other flavour observables [6].

One may ask what other processes are likely to be mediated by these couplings. At the individual coupling level, a comprehensive list can be found in [11]. At the product coupling level, all the $\lambda'\lambda'$ type products can mediate $B_d \to \nu\bar{\nu}$, which might be observable at the next generation B factories. Some of them can mediate four-quark interactions, like $\lambda'_{312}\lambda'_{332}$ mediating $b \to s\bar{s}d$ (or $B_s \to \phi K_S$), but the data is again inadequate to put further constraints. What we may emphasize is that these channels are worth looking into. If lucky, one might even hope for some LFV top decays too, like $t \to u\mu\tau$. In colliders like LHC, depending on the SUSY parameter space, one may observe a stau decaying into jets.

What happens if the RPV couplings are hierarchical not at the mass basis but at the weak or flavour basis? In that case, one has to rotate these couplings to the mass basis by some CKM-like mechanism, and this involves assumptions about the mixing matrices in the right-chiral quark sector. However, a general trend would be the appearance of the complete set of all RPV couplings at the mass basis. With the neutrino bounds at work, the constraints are expected to be much tighter, and hence less allowed parameter space for the charged Higgs. The constraints are more lenient if the mixing is in the up-quark sector, and our results do not change much. If the mixing is in the down-quark sector, the constraints are tighter, but one is still able to restore the low $\tan\beta$ region at least for some of the product couplings.

4. Summary and conclusions

The work was motivated by the fact that the tension between theory and experiment for the decay width of $B \to \tau\nu$ requires at least another tree-level contribution compatible in strength with the SM amplitude. The most plausible candidate is a charged Higgs; however, the contribution interferes destructively with the SM one, and one gets only a fine-tuned region where the solution exists. Moreover, even this region is highly disfavoured by other flavour observables.

The next option is to use a model where the 2HDM is embedded, like supersymmetry. The R-parity conserving version has some one-loop corrections to the $B \to \tau\nu$ amplitude, and that hardly helps alleviating the tension. On the other hand, if one invokes R-parity violation, there are more tree-level contributions to the decay, and the interference can be constructive. Thus, the tension on the charged Higgs parameter space is relieved, and one can have a massive charged Higgs at a sufficiently low value of $\tan\beta$. One can also rephrase the conclusion: if a supersymmetric charged Higgs is indeed found in this region, it will be worth-while to look for physics like R-parity violation, assuming the data on $B \to \tau\nu$ stands the test of time. As a consistency check, we have made sure that the RPV couplings satisfy all the existing bounds.

Acknowledgements

R.B. acknowledges the University Grants Commission, Government of India, for financial support. The work of A.K. was supported by CSIR, Government of India, and the DRS programme of the University Grants Commission.

References

[1] J. Charles, et al., arXiv:1106.4041 [hep-ph]; also V. Niess, CKMfitter Group, talk given at EPS 2011, Grenoble, France, July 2011.
[2] Several Lattice results, averaged by the CKMfitter group, and available at http://ckmfitter.in2p3.fr/www/results/plots_lp11/ckmEval_results_lp11.pdf.
[3] B. Aubert, et al., BaBar Collaboration, Phys. Rev. D 77 (2008) 011107(R);
I. Adachi, et al., Belle Collaboration, arXiv:0809.3834 [hep-ex];
B. Aubert, et al., BaBar Collaboration, Phys. Rev. D 81 (2010);
K. Hara, et al., Belle Collaboration, arXiv:1006.4201 [hep-ex];
P. del Amo Sanchez, et al., BaBar Collaboration, arXiv:1008.0104 [hep-ex].
[4] W.-S. Hou, Phys. Rev. D 48 (1993) 2342.
[5] A.G. Akeroyd, S. Recksiegel, J. Phys. G 29 (2003) 2311, arXiv:hep-ph/0306037.
[6] O. Deschamps, et al., Phys. Rev. D 82 (2010) 073012, arXiv:0907.5135 [hep-ph].
[7] G. Blankenburg, G. Isidori, arXiv:1107.1216 [hep-ph].
[8] A.G. Akeroyd, S. Recksiegel, Phys. Lett. B 541 (2002) 121.
[9] Y. Aida, E. Asakawa, G.-C. Cho, H. Matsuo, Phys. Rev. D 82 (2010) 115008.
[10] G.-C. Cho, H. Matsuo, arXiv:1107.3004 [hep-ph].
[11] R. Barbier, et al., Phys. Rep. 420 (2005) 1, arXiv:hep-ph/0406039.
[12] H.K. Dreiner, M. Kramer, B. O'Leary, Phys. Rev. D 75 (2007) 114016.
[13] J.P. Saha, A. Kundu, Phys. Rev. D 66 (2002) 054021.
[14] B.C. Allanach, A. Dedes, H.K. Dreiner, Phys. Rev. D 69 (2004) 115002;
B.C. Allanach, A. Dedes, H.K. Dreiner, Phys. Rev. D 72 (2005) 079902, Erratum.
[15] Y. Kao, T. Takeuchi, arXiv:0910.4980 [hep-ph].
[16] E. Komatsu, et al., WMAP Collaboration, Astrophys. J. Suppl. 192 (2011) 18.
[17] J.P. Saha, A. Kundu, Phys. Rev. D 70 (2004) 096002.