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New physics contribution to $B_s \to \mu^+\mu^-$ within R-parity violating supersymmetric models

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We revisit the problem of new physics (NP) contribution to the branching ratio of the $B_s \to \mu^+\mu^-$ decay in light of the recent observation of this decay by LHCb. We consider R-parity violating (RPV) supersymmetric models as a primary example—recently one has reported stringent constraints on the products of the RPV coupling constants that account for the $B_s \to \mu^+\mu^-$ transition at the tree level. We argue that despite the fact that the LHCb measurement of the $B(B_s \to \mu^+\mu^-)$ is in a remarkable agreement with the Standard Model (SM) prediction, there is still a room for a significant new physics contribution to the $B(B_s \to \mu^+\mu^-)$, as the sign of the $B_s \to \mu^+\mu^-$ transition amplitude may be opposite to that of the Standard Model; alternatively the amplitude may have a large phase. We conduct our analysis mainly for the case of real RPV couplings. We find that taking into account the scenario with the sign flip of the $B_s \to \mu^+\mu^-$ amplitude (as compared to that of the SM) makes the bounds on the RPV coupling products significantly weaker. Also, we discuss briefly how our results are modified if the RPV couplings have large phases. In particular, we examine the dependence of the derived bounds on the phase of the NP amplitude.

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The rare $B_s \to \mu^+\mu^-$ decay is believed to be one of the most powerful tools to test the physics that may occur beyond the Standard Model. Within the Standard Model this decay is loop induced and in addition is helicity suppressed. Numerical evaluation gives [1–3]

$$B(B_s \to \mu^+\mu^-) = (3.25 \pm 0.17) \times 10^{-9}. \tag{1}$$

In contrast, the $B_s \to \mu^+\mu^-$ decay rate may be dramatically enhanced within some of the Standard Model (SM) extensions and may exceed the SM prediction by several orders of magnitude. At the same time this decay is characterized by a pure final leptonic state, which causes the theoretical predictions for it to be very clean. It was therefore used intensively to constrain the SM extensions, and there was a hope to observe a distinct new physics (NP) signal in this decay mode.

Recently the LHCb Collaboration has reported the first evidence for the $B_s \to \mu^+\mu^-$ decay at 3.5σ level [4],

$$\bar{B}_{\rm exp}(B_s \to \mu^+\mu^-) = (3.2^{+1.5}_{-1.2}) \times 10^{-9}, \tag{2}$$

which is in a remarkable agreement with the SM prediction. However, it would not be correct to declare that there is no new physics contribution to $B_s \to \mu^+\mu^-$ at all. Some of the popular SM extensions do predict indeed a negligible NP contribution to the $B_s \to \mu^+\mu^-$ decay rate, due to strong correlations between the B_s-$\bar{B}_s$ mixing and $B_s \to \mu^+\mu^-$ amplitudes [5]. Yet, for other SM extensions the problem of new physics contribution to $B_s \to \mu^+\mu^-$ in light of the recent observation of this decay by LHCb is the subject of discussion in the literature [2,6–9]. In particular, it has been argued in [6–8] that the LHCb result still leaves room for a non-negligible NP contribution, due to the uncertainty in the experimental value of the $B(B_s \to \mu^+\mu^-)$.

In this paper we examine a source of new physics contribution to $B_s \to \mu^+\mu^-$ that would be actual even in the idealized limit of zero experimental and theoretical uncertainties in the $B(B_s \to \mu^+\mu^-)$ and perfect coincidence of the SM prediction with the experimental data. Namely, we consider a possibility for the $B_s \to \mu^+\mu^-$ transition amplitude to have a sign opposite to that of the Standard Model or to have a large phase. The LHCb measurement of the $B_s \to \mu^+\mu^-$ branching ratio constrains the decay rate, whereas the sign (if it is real) or the phase (if it is complex) of the transition amplitude remains arbitrary. Thus, it is possible that

(i) If the amplitude is real (or has a small enough phase so that it may be discarded), one may fit the experimental data for the $B(B_s \to \mu^+\mu^-)$ in particular when

$$\begin{aligned} &A^{\rm NP}(B_s \to \mu^+\mu^-) \\ &\quad \simeq -2A^{\rm SM}(B_s \to \mu^+\mu^-), \quad \text{so that} \\ &A(B_s \to \mu^+\mu^-) \\ &\quad = A^{\rm SM}(B_s \to \mu^+\mu^-) + A^{\rm NP}(B_s \to \mu^+\mu^-) \\ &\quad \simeq -A^{\rm SM}(B_s \to \mu^+\mu^-). \end{aligned} \tag{3}$$

(ii) If instead the NP amplitude has a large phase, one may fit the experimental data for the $B(B_s \to \mu^+\mu^-)$ when

$$\begin{aligned} &|A^{\rm SM}(B_s \to \mu^+\mu^-) + |A^{\rm NP}(B_s \to \mu^+\mu^-)|e^{i\Phi_{\rm NP}}| \\ &\quad \simeq |A^{\rm SM}(B_s \to \mu^+\mu^-)| \end{aligned} \tag{4}$$

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[if neglecting the SM amplitude phase and using the approximation $B_{\rm exp}(B_s \to \mu^+\mu^-) \approx B^{\rm SM}(B_s \to \mu^+\mu^-)$]. Note that Eq. (4) implies

$$-2A^{\rm SM}(B_s \to \mu^+\mu^-) < {\rm Re}[A^{\rm NP}(B_s \to \mu^+\mu^-)] < 0, \tag{5}$$

$$|{\rm Im}[A^{\rm NP}(B_s \to \mu^+\mu^-)]| \lesssim |A^{\rm SM}(B_s \to \mu^+\mu^-)|. \tag{6}$$

In particular,

$$\begin{aligned} &|{\rm Im}[A^{\rm NP}(B_s \to \mu^+\mu^-)]| \\ &\quad \simeq |A^{\rm SM}(B_s \to \mu^+\mu^-)| \\ &\text{if } {\rm Re}[A^{\rm NP}(B_s \to \mu^+\mu^-)] \\ &\quad \simeq -A^{\rm SM}(B_s \to \mu^+\mu^-). \end{aligned} \tag{7}$$

One may infer from Eqs. (3)–(7) that the NP contribution to the $B_s \to \mu^+\mu^-$ transition amplitude is the largest when the amplitude just flips the sign as compared to that of the Standard Model (rather than getting a large nontrivial phase). So, we will be concentrating here mainly on the case of a real amplitude, by assuming that the relevant NP parameters are real. We will however discuss at the end of the paper how our results are modified in presence of large phases of the NP parameters.

Note that the possibility of the $B_s \to \mu^+\mu^-$ amplitude sign flip has already been mentioned in [7] where one considered new physics models with modified Z-boson couplings to down-type quarks. This possibility has been rejected there, as it is disfavored by the constraints on $Z \to b\bar{b}$. To our best knowledge, there is no reason to disfavor the $B_s \to \mu^+\mu^-$ amplitude sign flip within other SM extensions (in fact it has also been implicitly considered in [2] within the general analysis of the NP contribution to $B_s \to \mu^+\mu^-$ in a variety of models, with the amplitude phases varied freely from 0 to π). In our opinion, the detailed analysis of the possibility that the $B_s \to \mu^+\mu^-$ amplitude may have a sign opposite to that of the Standard Model (or have a large phase) may be of great importance, especially in light of future improvement of the experimental accuracy of measurements of the $B_s \to \mu^+\mu^-$ branching ratio.

We consider here R-parity violating supersymmetric models with leptonic number violation as a primary example. It has been recently argued [9] that the remarkable agreement between the LHCb measurement and the SM prediction for the $B_s \to \mu^+\mu^-$ branching ratio implies rigorous constraints on the RPV coupling products that account for the $B_s \to \mu^+\mu^-$ transition at the tree level. We show that if the $B_s \to \mu^+\mu^-$ transition amplitude is allowed to have a sign opposite to that of the Standard Model, bounds on the RPV couplings may be by order of magnitude weaker.

The most general Yukawa superpotential for an explicitly broken R-parity supersymmetric theory may be written as

$$W_{\not R} = \frac{1}{2}\lambda_{ijk}L_iL_jE_k^c + \lambda'_{ijk}L_iQ_jD_k^c + \frac{1}{2}\lambda''_{ijk}U_i^cD_j^cD_k^c. \tag{8}$$

Here Q and L denote $SU(2)_L$ doublet quark and lepton superfields, and U, D and E stand for $SU(2)_L$ singlet up-quark, down-quark and charged lepton superfields. Also, i, j, $k = 1, 2, 3$ are generation indices. We shall require baryon number symmetry by setting λ''_{ijk} to zero. Also, as mentioned above, we will assume the couplings λ_{ijk} and λ'_{ijk} are real.

Subsequently, the Lagrangian describing the RPV supersymmetry (SUSY) contribution to $B_s \to \mu^+\mu^-$ can be written as

$$\begin{aligned} \mathcal{L}_{\not R} = -(&\lambda'_{i23}\tilde{\nu}_{i_L}\bar{b}P_Ls + \lambda'_{i32}\tilde{\nu}_{i_L}\bar{s}P_Lb + \lambda_{i22}\tilde{\nu}_{i_L}\bar{\mu}P_L\mu \\ &+ \lambda'_{2k2}\tilde{u}_{k_L}\bar{s}P_L\mu + \lambda'_{2k3}\tilde{u}_{k_L}\bar{b}P_L\mu + {\rm H.c.}) \end{aligned} \tag{9}$$

where $P_{L,R}$ are the helicity projection operators, and we use the notation $P_L = (1-\gamma_5)/2$. Note that for the sake of transparency of our analysis, we neglect the transformation of the RPV couplings from the weak isospin basis to the (s) quark and sneutrino mass basis. (We invoke however to the reader to be cautious when using the bounds on RPV coupling products derived in this paper. Rigorously speaking, they may be used for the processes involving down-type quark–down type quark–sneutrino and down-type quark–up-type squark–charged lepton transitions only.)

Within R-parity violating supersymmetric models, to the lowest order in perturbation theory the $B_s \to \mu^+\mu^-$ transition occurs at the tree level, due to exchange of sneutrinos or up-type squarks, as depicted in Fig. 1. We need also to include the SM contribution to $B_s \to \mu^+\mu^-$: recall that we are interested in destructive interference of the SM and NP amplitudes. Thus, the relevant low-energy $|\Delta B| = 1$ effective Hamiltonian would have the following form:

$$H_{\rm eff}^{\Delta B=1} = H_{\rm eff}^{\rm SM} + H_{\rm eff}^{\tilde{\nu}} + H_{\rm eff}^{\tilde{u}}. \tag{10}$$

Here [10,11]

$$H_{\rm eff}^{\rm SM} = \frac{-4G_F}{\sqrt{2}}\frac{\alpha}{2\pi\sin^2\theta_W} \times (V_{tb}^\star V_{ts})\eta_Y Y_0(x_t)\bar{b}\gamma^\nu P_L s\bar{\mu}\gamma_\nu P_L\mu + {\rm H.c.} \tag{11}$$

where [12]

$$Y_0(x_t) = \frac{x_t}{8}\left(\frac{4-x_t}{1-x_t} + \frac{3x_t}{(1-x_t)^2}\ln x_t\right),$$

$x_t = m_t^2/M_W^2$, and η_Y is the factor that accounts for the QCD corrections to $Y_0(x_t)$.

Two other terms in Eq. (10) are derived by integrating out the sneutrino and squark heavy degrees of freedom.

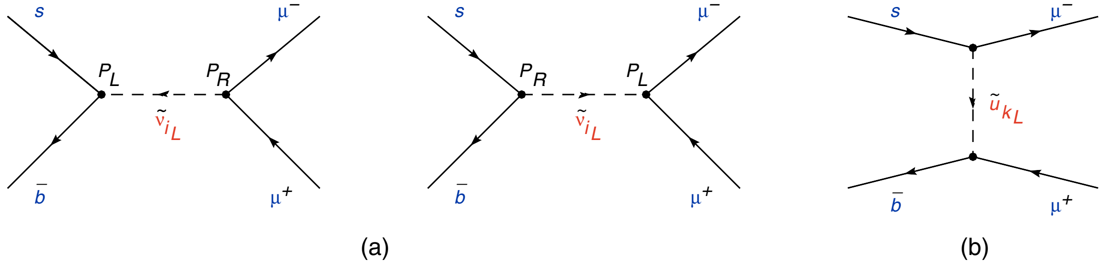

FIG. 1 (color online). Diagrams for the $B_s \to \mu^+\mu^-$ transition within R-parity violating supersymmetric models to the lowest order in perturbation theory, (a) due to exchange of sneutrinos, (b) due to exchange of up-type squarks. The direction of the sneutrino propagator depends on the helicities of the quark and lepton states, in other words whether we have P_L or we have P_R operator at an interaction vertex.

This yields

$$H_{\rm eff}^{\tilde\nu} = -\left(\frac{\lambda^\star_{i22}\lambda'_{i23}}{m^2_{\tilde\nu_{i_L}}}\bar b P_L s\bar\mu P_R\mu + \frac{\lambda'^\star_{i32}\lambda_{i22}}{m^2_{\tilde\nu_{i_L}}}\bar b P_R s\bar\mu P_L\mu + \text{H.c.}\right) \tag{12}$$

$$H_{\rm eff}^{\tilde u} = \frac{\lambda'^\star_{2k2}\lambda'_{2k3}}{2m^2_{\tilde u_{k_L}}}\bar b\gamma^\nu P_R s\bar\mu\gamma_\nu P_L\mu + \text{H.c.} \tag{13}$$

where $\tilde\nu_{i_L}$, $\tilde u_{k_L}$ are respectively the lightest sneutrino and the lightest "left" up-type squark states.[1] For (nearly) degenerate sneutrino and/or squark masses, one should replace in Eq. (12) and/or Eq. (13) the lightest sparticle masses by universal sneutrino and/or squark masses, $m_{\tilde\nu_{i_L}} \to m_{\tilde\nu_L}$, $m_{\tilde u_{k_L}} \to m_{\tilde u_L}$, as well as sum over indices i and/or k.

Also, for the sake of clarity of our analysis, we will follow Ref. [5] and assume

$$\lambda'_{i23} = \lambda'_{i32} \tag{14}$$

in our further calculations.

Using Eqs. (10)–(13) as well as the simplifying assumption (14), one may present the $B_s \to \mu^+\mu^-$ transition amplitude in the following form:

$$A(B_s\to\mu^+\mu^-) = A^{\rm SM}(B_s\to\mu^+\mu^-) + A^{\tilde\nu}(B_s\to\mu^+\mu^-) + A^{\tilde u}(B_s\to\mu^+\mu^-) \tag{15}$$

where

$$A^{\rm SM}(B_s\to\mu^+\mu^-) = -\langle\mu^+\mu^-|H_{\rm eff}^{\rm SM}|B_s\rangle = \frac{-iG_F}{\sqrt2}\frac{\alpha f_{B_s}m_\mu}{\pi\sin^2\theta_W}(V^\star_{tb}V_{ts})\eta_Y Y_0(x_t)\bar u(p_-)\gamma_5 v(p_+) \tag{16}$$

$$A^{\tilde\nu}(B_s\to\mu^+\mu^-) = -\langle\mu^+\mu^-|H_{\rm eff}^{\tilde\nu}|B_s\rangle = \frac{-i\lambda^\star_{i22}\lambda'_{i23}f_{B_s}M^2_{B_s}}{2m^2_{\tilde\nu_{i_L}}m_b}\bar u(p_-)\gamma_5 v(p_+) \tag{17}$$

$$A^{\tilde u}(B_s\to\mu^+\mu^-) = -\langle\mu^+\mu^-|H_{\rm eff}^{\tilde u}|B_s\rangle = \frac{-i\lambda'^\star_{2k2}\lambda'_{2k3}f_{B_s}m_\mu}{4m^2_{\tilde u_{k_L}}}\bar u(p_-)\gamma_5 v(p_+) \tag{18}$$

where $u(p_-)$ and $v(p_+)$ are the bispinor wave functions of the leptonic states. (Subsequently, p_+ and p_- are the momenta of μ^+ and μ^-.) In deriving (16)–(18) we used the following parametrization of the hadronic matrix elements:

$$\langle 0|\bar b\gamma^\nu P_L s|B_s\rangle = \frac{if_{B_s}}{2}p_B^\nu$$
$$\langle 0|\bar b\gamma^\nu P_R s|B_s\rangle = \frac{-if_{B_s}}{2}p_B^\nu$$
$$\langle 0|\bar b P_R s|B_s\rangle = \frac{if_{B_s}}{2}\left(\frac{M^2_{B_s}}{m_b}\right)$$
$$\langle 0|\bar b P_L s|B_s\rangle = \frac{-if_{B_s}}{2}\left(\frac{M^2_{B_s}}{m_b}\right)$$

where f_{B_s} is the B_s meson decay constant, and p_B is the B_s 4-momentum.

We want to stress that all three parts of the amplitude have the same structure. They all contain the same pseudoscalar bispinor bilinear form multiplied by some factor. In what follows, both $A^{\tilde\nu}(B_s\to\mu^+\mu^-)$ and $A^{\tilde u}(B_s\to\mu^+\mu^-)$ may interfere with the SM amplitude. In other

[1] It is assumed that squark mass eigenstates do not differ significantly from the left and right squark states. This is known to be the case for most SUSY scenarios with the squark masses much greater than 100 GeV.

words, the $B_s \to \mu^+\mu^-$ amplitude may have a sign opposite to that of the SM both due to the contribution of the sneutrino-mediated diagrams, and due to the contribution of the squark-mediated diagram.[2]

Calculation of the decay branching ratio using (15)–(18) is straightforward and yields

$$B(B_s \to \mu^+\mu^-) = \frac{\tau_{B_s} M_{B_s} f_{B_s}^2}{8\pi}\sqrt{1-\frac{4m_\mu^2}{M_{B_s}^2}}\left|\frac{G_F}{\sqrt{2}}\frac{\alpha m_\mu}{\pi \sin^2\theta_W} \times (V_{tb}^\star V_{ts})\eta_Y Y_0(x_t) + \frac{\lambda_{i22}^\star \lambda'_{i23} M_{B_s}^2}{2m_{\tilde{\nu}_{i_L}}^2 m_b} + \frac{\lambda_{2k2}^{\prime\star}\lambda'_{2k3} m_\mu}{4m_{\tilde{u}_{k_L}}^2}\right|^2 \tag{19}$$

where τ_{B_s} is the average lifetime of the B_s meson.

We use during the numerical analysis $\tau_{B_s} = 1.509$ ps [13], $f_{B_s} = 0.225$ GeV [14], $V_{tb}^\star V_{ts} = 0.0405$ [15] (as mentioned above, we neglect the small phase of this CKM product), $M_W = 80.4$ GeV, $\sin^2\theta_W = 0.231$, $G_F = 1.166 \times 10^{-5}$ GeV^{-2}, $\alpha = \alpha(M_Z) = 1/128$, $m_\mu = 0.106$ GeV, $M_{B_s} = 5.3667$ GeV, $m_b = \overline{m}_b(m_b) = 4.18$ GeV [16]. For the top quark mass we use $m_t^{\rm pole} = 173.2$ GeV [17], which yields for the $\overline{\rm MS}$, QCD renormalized mass $\bar{m}_t(m_t) = 163.2$ GeV [1]; $\eta_Y = 1.012$ for $x_t = \bar{m}_t^2(m_t)/M_W^2$ [1,18–20].

We neglect the uncertainties in the values of the input parameters specified above. Those are known to alter the predictions for the $B(B_s \to \mu^+\mu^-)$ by about 10% [1,2]. This uncertainty in the $B(B_s \to \mu^+\mu^-)$ is much less than the one in the experimental value of the branching ratio and the one in our results due to destructive interference of different NP amplitudes (see the discussion at the end of the paper).

We choose $m_{\tilde{\nu}_{i_L}} \gtrsim 100$ GeV and $m_{\tilde{u}_{k_L}} \gtrsim 500$ GeV. The squark masses below 500 GeV are highly disfavored by the LHC data (see [16,21] and references therein). To our best knowledge, however, no such strong constraints on sneutrino masses has been derived so far [16].

Also, following the common approach, we will assume only one nonvanishing RPV coupling product at a time, or alternatively only one of the NP amplitudes in (15) to be nonvanishing at a time.

We consider first an idealized scenario with zero uncertainties in the experimental and theoretical values of the $B_s \to \mu^+\mu^-$ branching ratio and perfect coincidence of the Standard Model prediction with the experimental data. In such a scenario (if assuming real RPV couplings), nonvanishing new physics contribution to $B_s \to \mu^+\mu^-$ may occur if only the transition amplitude has a sign opposite to that of the Standard Model. Following the approach of one nonvanishing coupling product at a time, we choose first $\lambda_{2k2}^{\prime\star}\lambda'_{2k3} = 0$ or equivalently $A^{\tilde{u}}(B_s \to \mu^+\mu^-) = 0$. Then the transition amplitude flips the sign if

$$A^{\tilde{\nu}}(B_s \to \mu^+\mu^-) = -2A^{\rm SM}(B_s \to \mu^+\mu^-).$$

Using Eqs. (16) and (17) and the values of the input parameters specified above, one finds that this occurs when

$$-\lambda_{i22}^\star \lambda'_{i23} = 2.12 \times 10^{-6}\left(\frac{m_{\tilde{\nu}_{i_L}}}{100 \text{ GeV}}\right)^2. \tag{20}$$

This value of $\lambda_{i22}^\star\lambda'_{i23}$ is several times greater in magnitude than the bound quoted in [9] (as no amplitude sign flip or large phase has been considered in [9]). Nevertheless, Eq. (20) implies rigorous constraints on this coupling product or alternatively on the sneutrino masses. Indeed, Eq. (20) implies $(-\lambda_{i22}^\star\lambda'_{i23}) \sim 10^{-6}$ for the lightest sneutrino mass ~100 GeV. Alternatively, if one desires for this coupling product to be of the same order as the SM weak coupling squared ($g^2 \sim 0.5$), the lightest sneutrino should have a mass ~50 TeV. This is a manifestation of the so-called flavor problem [22,23]: to assure that tree level flavor changing neutral currents beyond the SM do not conflict with the experimental data, either the relevant couplings should be unnaturally small or the new physics mass scale should be enormously large. Solving the flavor problem goes beyond the scope of the present paper. Instead we will simply assume further that $\lambda_{i22}^\star\lambda'_{i23} = 0$, or $A^{\tilde{\nu}}(B_s \to \mu^+\mu^-)$ vanishes, and we will be concentrating on the contribution of the squark-mediated diagram only [Fig. 1(b)]. As mentioned above, possible effects of interference of different NP amplitudes will be discussed at the end of the paper.

If assuming $A^{\tilde{\nu}}(B_s \to \mu^+\mu^-) = 0$, the transition amplitude flips the sign when

$$A^{\tilde{u}}(B_s \to \mu^+\mu^-) = -2A^{\rm SM}(B_s \to \mu^+\mu^-).$$

Using Eqs. (16) and (18) and the values of the input parameters specified above, one finds that this occurs when

$$-\lambda_{2k2}^{\prime\star}\lambda'_{2k3} = 6.88 \times 10^{-3}\left(\frac{m_{\tilde{u}_{k_L}}}{500 \text{ GeV}}\right)^2. \tag{21}$$

Equation (21) implies rather weak constraints on the couplings λ'_{2k2} and λ'_{2k3}. If assuming no hierarchy in the values of λ'_{2k2} and λ'_{2k3}, one gets $|\lambda'_{2k2}| \sim 0.085$ and $|\lambda'_{2k3}| \sim 0.085$ for $m_{\tilde{u}_{k_L}} \sim 500$ GeV. Thus, moderately small values of λ'_{2k2} and λ'_{2k3} are still allowed for $m_{\tilde{u}_{k_L}} \sim 500$ GeV. Furthermore, choosing the lightest left up-type squark mass to be heavier (say 1 TeV of few TeV) would yield larger values for λ'_{2k2} and λ'_{2k3} (and for their product) to be allowed.

At first glance this result is not surprising, as the contribution of the diagram with a squark exchange in Fig. 1(b) is helicity suppressed (like the SM contribution),

[2]If we give up the simplifying assumption (14), $A^{\tilde{\nu}}(B_s \to \mu^+\mu^-)$ will also contain a term with a scalar bispinor bilinear form. This term however will not interfere with the other terms of the transition amplitude, so it does not play any essential role in our analysis.

as can be seen e.g. from Eq. (18). We want to stress however that this is a rather nontrivial result, in a sense that one should consider the possibility of the $B_s \to \mu^+\mu^-$ amplitude sign flip to derive it. If instead one assumes that the sign of the transition amplitude is the same as within the SM, so that the NP contribution is solely due to the uncertainty in the experimental value of the $B(B_s \to \mu^+\mu^-)$, the constraints on the coupling product $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ are significantly stronger. To illustrate this, we will consider a realistic scenario now: we will demand that our predictions for the $B_s \to \mu^+\mu^-$ branching ratio fall in the experimentally allowed interval.

In order to do this, one should take into account that the experimentally measured branching ratio of the $B_s \to \mu^+\mu^-$ decay is the time integrated branching ratio [usually denoted $\bar{B}(B_s \to \mu^+\mu^-)$ like in Eq. (2) above]. It is related to the "theoretical" branching ratio as [2,24–26]

$$B(B_s \to \mu^+\mu^-) = \left(\frac{1-y_s^2}{1+A^{\mu\mu}_{\Delta\Gamma}y_s}\right)\bar{B}(B_s \to \mu^+\mu^-). \tag{22}$$

Here [27]

$$y_s = \frac{\Delta\Gamma_s}{2\Gamma_s} = 0.088 \pm 0.014 \tag{23}$$

where $\Delta\Gamma_s$ is the width difference in the B_s-$\bar{B}_s$ mixing, and Γ_s is the average width of the B_s meson. The expression for $A^{\mu\mu}_{\Delta\Gamma}$ in terms of Wilson coefficients of the low-energy effective operators may be found in [2]. For the considered case of real NP couplings and under simplifying assumption (14), one can show after doing some algebra that $A^{\mu\mu}_{\Delta\Gamma} = 1$. Thus, the experimentally allowed (1σ) interval for the $\bar{B}(B_s \to \mu^+\mu^-)$ [given by Eq. (2)] is converted to the following allowed interval for the theoretical branching ratio:

$$\begin{aligned} B(B_s \to \mu^+\mu^-) &= (1-y_s)\bar{B}(B_s \to \mu^+\mu^-) \\ &= (2.9^{+1.4}_{-1.2}) \times 10^{-9}. \end{aligned} \tag{24}$$

Equation (24) [combined with Eq. (19) in the limit when only the squark-mediated diagram in Fig. 1(b) gives a nonvanishing NP contribution] yields the following constraints on the coupling product $\lambda'^{\star}_{2k2}\lambda'_{2k3}$:

$$\begin{aligned} &-4.9 \times 10^{-4}\left(\frac{m_{\tilde{u}_{k_L}}}{500\ \mathrm{GeV}}\right)^2 \\ &\quad \le -\lambda'^{\star}_{2k2}\lambda'_{2k3} \le 9.6 \times 10^{-4}\left(\frac{m_{\tilde{u}_{k_L}}}{500\ \mathrm{GeV}}\right)^2 \end{aligned} \tag{25}$$

and

$$\begin{aligned} &5.92 \times 10^{-3}\left(\frac{m_{\tilde{u}_{k_L}}}{500\ \mathrm{GeV}}\right)^2 \\ &\quad \le -\lambda'^{\star}_{2k2}\lambda'_{2k3} \le 7.37 \times 10^{-3}\left(\frac{m_{\tilde{u}_{k_L}}}{500\ \mathrm{GeV}}\right)^2. \end{aligned} \tag{26}$$

The first interval [given by (25)] is derived when the $B_s \to \mu^+\mu^-$ transition amplitude has the same sign as that of the Standard Model. The new physics contribution is due to the uncertainty in the experimental value of the $B_s \to \mu^+\mu^-$ branching ratio. This interval for $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ is in a reasonable agreement with that quoted in Ref. [9]. The second interval [given by (26)] is derived when the transition amplitude has a sign opposite to that of the Standard Model. In that case the allowed values of $-\lambda'^{\star}_{2k2}\lambda'_{2k3}$ are greater by an order of magnitude. As discussed above, this implies weaker constraints on the allowed region of the NP parameter space.

Notice also that for the $B_s \to \mu^+\mu^-$ amplitude to flip the sign, the coupling product $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ must be negative [as it follows from Eq. (26)]. Contrary to this, within the other interval [given by (25)], the sign of $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ is arbitrary.

We used the 1σ experimental interval to derive the constraints on $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ given by (25) and (26). A more conservative approach would imply using the 95% C.L. interval, $\bar{B}(B_s \to \mu^+\mu^-) = [1.1 \div 6.4] \times 10^{-9}$ [4]. One would observe the same effect in that case as well, although less pronounced and harder to analyze. While using the 95% C.L. interval (instead of the 1σ one) would affect the sign-flip interval [given by (26)] by about 10% only, the same-sign interval would be significantly more widespread than (25). We leave for a reader to verify that if using the 95% C.L. interval, the maximum value of $-\lambda'^{\star}_{2k2}\lambda'_{2k3}$ in the sign-flip interval would be about 5 times greater than the maximum value of $|\lambda'^{\star}_{2k2}\lambda'_{2k3}|$ in the same-sign interval, or constraints on this coupling product would still be significantly weaker when taking into account the possibility of the $B_s \to \mu^+\mu^-$ amplitude sign flip.

In principle, one may conduct a similar analysis for the contribution of the sneutrino-mediated diagrams on Fig. 1(a) and subsequently for the other coupling product, $\lambda^{\star}_{i22}\lambda'_{i23}$. Assuming now that the squark-mediated diagram in Fig. 1(b) has a vanishing contribution to $B_s \to \mu^+\mu^-$, one will get in this case two different intervals for $\lambda^{\star}_{i22}\lambda'_{i23}$ [that originate in the same way as (25) and (26) for $\lambda'^{\star}_{2k2}\lambda'_{2k3}$]. We leave this for a reader as another exercise to do.

In a realistic scenario neither of the diagrams in Fig. 1 may have a vanishing contribution to $B_s \to \mu^+\mu^-$. In addition, one should also take into account the impact of the R-conserving sector of the theory on the $B_s \to \mu^+\mu^-$ transition amplitude as well [6]. Thus, in a realistic scenario one has different sources of a NP contribution to $B_s \to \mu^+\mu^-$ that may in general interfere both constructively and destructively [28].

If different NP amplitudes interfere constructively, the coupling product $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ may also acquire the values between the two intervals given by (25) and (26). [That is to say, the $B_s \to \mu^+\mu^-$ amplitude sign flip may be only in part due to the contribution of the squark-mediated diagram in Fig. 1(b); it may also be in part due to other new physics effects.] In other words, one should replace (25) and (26) by

$$-4.9\times 10^{-4}\left(\frac{m_{\tilde{u}_{k_L}}}{500\ \mathrm{GeV}}\right)^2 \leq -\lambda'^{\star}_{2k2}\lambda'_{2k3} \leq 7.37\times 10^{-3}\left(\frac{m_{\tilde{u}_{k_L}}}{500\ \mathrm{GeV}}\right)^2. \quad (27)$$

Of course, the NP amplitudes may also interfere destructively. In that case the bounds on $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ given by Eq. (27) may somehow be distorted (they may become weaker). Yet, if there is no fine-tuning or exact cancellation of the contributions of different NP amplitudes, it is very unlikely that this distortion alter the bounds on $\lambda'^{\star}_{2k2}\lambda'_{2k3}$, say, by an order of magnitude. Thus, one may always use (27) to get an insight into how large (in order of magnitude) the coupling product $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ is still allowed to be.

So far we were assuming that the R-parity violating couplings are real (or have small enough phases so that they may be discarded). Yet, our analysis may be extended also to the case when these couplings have large phases. Demanding again that our predictions fall into the experimentally allowed interval [and assuming again $A^{\tilde{u}}(B_s \to \mu^+\mu^-)$ to be the only nonvanishing NP amplitude], one may derive an upper bound on the absolute value of the coupling product $\lambda'^{\star}_{2k2}\lambda'_{2k3}$ as a function of the NP amplitude phase $\Phi_{\mathrm{NP}} = \arg(\lambda'^{\star}_{2k2}\lambda'_{2k3})$. One must however be cautious what the allowed interval is now, as the observable $A^{\mu\mu}_{\Delta\Gamma}$ is not equal to unity anymore. Thus, Eq. (24), which we were using in the case of real NP couplings, is not valid here. One shall demand instead

$$\left(\frac{1 + A^{\mu\mu}_{\Delta\Gamma} y_s}{1 - y_s^2}\right) B(B_s \to \mu^+\mu^-) = \bar{B}(B_s \to \mu^+\mu^-) = (3.2^{+1.5}_{-1.2})\times 10^{-9} \quad (28)$$

where in the limit of vanishing contribution of scalar operators, $A^{\mu\mu}_{\Delta\Gamma} = \cos(2\varphi_P - \phi_s^{\mathrm{NP}})$ [2]. Here ϕ_s^{NP} is the NP piece of the B_s-$\bar{B}_s$ mixing phase, and within the considered scenario φ_P is the phase of the total (SM + NP) amplitude. There is a rather weak correlation between the NP contribution to B_s-$\bar{B}_s$ coming from the R-parity violating sector and that to $B_s \to \mu^+\mu^-$ [5]. Moreover, this correlation is negligible, if analyzing the contribution of the squark-mediated diagram in Fig. 1(b) only (or analyzing the constraints on $\lambda'^{\star}_{2k2}\lambda'_{2k3}$). Also, the recent measurements of ϕ_s at LHCb [27], combined with the knowledge of the SM piece of ϕ_s, allow us to infer that $\phi_s^{\mathrm{NP}} \lesssim 0.15$ radians, so this phase is too small to affect $A^{\mu\mu}_{\Delta\Gamma}$ significantly. We will discard ϕ_s^{NP} in our calculations, thus using $A^{\mu\mu}_{\Delta\Gamma} \approx \cos 2\varphi_P$. Note that φ_P does not acquire a unique value as the NP amplitude phase Φ_{NP} is fixed. φ_P depends both on $\Phi_{\mathrm{NP}} = \arg(\lambda'^{\star}_{2k2}\lambda'_{2k3})$ (which is the only genuine free parameter in our analysis), and on $|\lambda'^{\star}_{2k2}\lambda'_{2k3}|/m^2_{\tilde{u}_{k_L}}$ (or on the relative weight of the NP amplitude compared to the SM one).

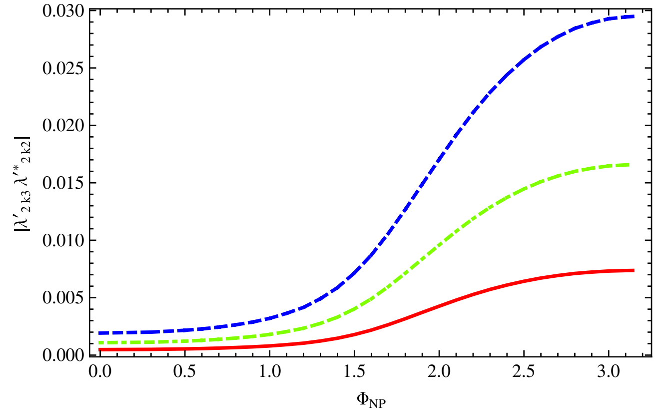

FIG. 2 (color online). Upper bound on $|\lambda'^{\star}_{2k2}\lambda'_{2k3}|$ as a function of the NP amplitude phase Φ_{NP} for $m_{\tilde{u}_{k_L}} = 500$ GeV (solid red line), $m_{\tilde{u}_{k_L}} = 750$ GeV (dashed-dotted green line), $m_{\tilde{u}_{k_L}} = 1$ TeV (dashed blue line).

The derived bound on $|\lambda'^{\star}_{2k2}\lambda'_{2k3}|$ as a function of the NP amplitude phase Φ_{NP} is presented in Fig. 2. As one can see from Fig. 2, the bound on $|\lambda'^{\star}_{2k2}\lambda'_{2k3}|$ becomes weaker as Φ_{NP} gets larger, and it is the weakest when $\Phi_{\mathrm{NP}} \to \pi$. As mentioned above, this result could also be inferred from the analysis of Eqs. (3)–(7). Thus, the most general bound on $|\lambda'^{\star}_{2k2}\lambda'_{2k3}|$ (both in the case when this product is real and in the case this product has a phase) would be

$$|\lambda'^{\star}_{2k2}\lambda'_{2k3}| \lesssim 7.37\times 10^{-3}\left(\frac{m_{\tilde{u}_{k_L}}}{500\ \mathrm{GeV}}\right)^2. \quad (29)$$

In conclusion, we have revisited the problem of new physics contribution to the $B_s \to \mu^+\mu^-$ decay in light of the recent experimental measurement of this decay branching ratio by the LHCb Collaboration. We have examined R-parity violating supersymmetric models as a primary example, and argued that there is still room for a significant NP contribution, as the transition amplitude still may have a sign opposite to that of the Standard Model or alternatively may get a large phase. We have found that if taking into account the effect of the $B_s \to \mu^+\mu^-$ amplitude possible sign flip as compared to that of the SM (or possible large phase), the bounds imposed on the RPV coupling products that account for the $B_s \to \mu^+\mu^-$ transition may be weaker by an order of magnitude than if the effect of the amplitude sign flip (or possible large phase) is disregarded. We emphasize that a similar effect may be observed also within other SM extensions. So considering within other new physics models the possibility for the $B_s \to \mu^+\mu^-$ transition amplitude to have a sign opposite to that of the SM or to have a large nontrivial phase is strongly encouraged.

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