Revisiting the $B$-physics anomalies in $R$-parity violating MSSM

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In recent years, several deviations from the Standard Model predictions in semileptonic decays of $B$-meson might suggest the existence of new physics which would break the lepton-flavour universality. In this work, we have explored the possibility of using muon sneutrinos and right-handed sbottoms to solve these $B$-physics anomalies simultaneously in $R$-parity violating minimal supersymmetric standard model. We find that the photonic penguin induced by exchanging sneutrino can provide sizable lepton flavour universal contribution for the $R_{B^0 \to \phi \mu \mu}$ transition. The photonic penguin amplitude is shown to be sizable and gives a $3.3\sigma$ discrepancy. However, the SM prediction is around 1 with $R_{B^0 \to \phi \mu \mu}$ decay. The SM prediction is around 1 with $R_{B^0 \to \phi \mu \mu}$ decay. However, there is a $2.6\sigma$ discrepancy in the low $q^2$ region and a $3.0\sigma$ discrepancy in the high $q^2$ region, respectively. The Belle collaboration also reported their measurements of $R_{K^*}$ [5, 6], which are consistent with the SM predictions within their quite large error bars. In addition to $R_{K^*}$, there are also some other deviations in $b \to s\mu^+\mu^-$ transition, such as the angular observables $\alpha_{K^+}$ [7–9] of $B \to K^+ \mu^+\mu^-$ decay with $2.6\sigma$ discrepancy [10–15] and the differential branching fraction of $B_s \to \phi \mu^+\mu^-$ decay with $3.3\sigma$ discrepancy [16, 17].

These deviations indicate the possible existence of new physics (NP) beyond the SM in $b \to s\ell^+\ell^-$ transition. This NP may break LFU. Many recent model-independent analyses [18–25] show that some scenarios can explain the $b \to s\ell^+\ell^-$ anomaly well. To express the fit results, we consider the low-energy effective weak Lagrangian governing the $b \to s\ell^+\ell^-$ transition

$$\mathcal{L}_{\text{eff}} = \frac{4G_F}{\sqrt{2}} \eta \sum_i C_i O_i + \text{H.c.},$$

where CKM factor $\eta \ell = V_{\ell b} V_{\ell s}^*$. We mainly concern the semileptonic operators

$$O_9 = \frac{e^2}{16\pi^2} (\bar{s}_\gamma \mu \mu \ell), \quad O_{10} = \frac{e^2}{16\pi^2} (\bar{s}_\gamma \mu \mu \gamma_5 \ell),$$

while $P_L = (1 - \gamma_5)/2$ is the left-handed chirality projector. The Wilson coefficients $C_{9,10} = C_{9,10}^{\text{SM}} + C_{9,10}^{\text{NP}}$. In this work, we try to explain the anomaly through a two-parameter scenario where the total NP effects are given by [26]

$$C_{9,\mu,\tau} = C_9^{\text{NP}} + C_9^{\text{SM}}, \quad C_{10,\mu,\tau} = -C_9^{\text{NP}}, \quad C_{9,\tau,\mu} = C_9^{\text{NP}} = 0.$$
The combined measurements of $R(D^*)$ and $R(D)$ are from BaBar [28, 29] and Belle [30, 31], and Belle [32, 33] and LHCb [34–36] only give the measurements of $R(D^*)$. After being averaged by the Heavy Flavor Averaging Group (HFLAV) [37], they give the results as follows [38]

$$R(D)_{\text{avg}} = 0.340 \pm 0.027 \pm 0.013, \quad (9)$$

$$R(D^*)_{\text{avg}} = 0.295 \pm 0.011 \pm 0.008, \quad (10)$$

with a correlation of $-0.38$. Comparing these with the arithmetic average of the SM predictions [38–42],

$$R(D)_{\text{SM}} = 0.299 \pm 0.003, \quad R(D^*)_{\text{SM}} = 0.258 \pm 0.005, \quad (11)$$

one can see that the difference between experiment and theory is at about 3.08σ, implying the existence of LFU violating NP in the charged-current $B$-decays. Global analyses [43–47] show that the NP contributing to the left-handed operator $B_{\text{NP}}$ in the charged-current processes in the framework of MSSM [27] and the colour indices are suppressed. Based on this inspiration from the paper by Bauer and Neubert [61], the authors in Ref. [58] used sneutrinos to explain it and found that it is allowed by the Lagrangian which can be obtained by the chiral superfields composing of the fermions and sfermions as follows

$$\mathcal{L} = \lambda'_{ijk} \bar{L}_i d_j R_k + \bar{d}_j R_k L_i + d_k^{\ast} \bar{c}_l \bar{L}_j \bar{u}_l + H.c., \quad (13)$$

where the sparticles are denoted by “*”, and “c” indicates charge conjugated fields. Working in the mass eigenstates for the down type quarks and assuming sfermions are in their mass eigenstates, one replaces $u_{Lj}$ by $(V^{\dagger} u_{Lj})$ in Eq. (13).

These $R$-parity violating interactions can induce $b \rightarrow s \ell^+ \ell^-$ processes by exchanging left-handed up squarks $\tilde{u}_{Lj}$ at tree level, but resulting in the operators with right-handed quark current, which are unable to explain the $b \rightarrow s \ell^+ \ell^-$ anomaly. This unwanted effect can be eliminated by assuming that the masses of $\tilde{u}_{Lj}$ are very large or/and by assuming that $\lambda'_{ij2} = 0$. Assuming that $\lambda_{ij2} = 0$ also forbids the exchange of $\tilde{L}_i$ or/and $\tilde{d}_{Lj}$ in one loop level to affect the $b \rightarrow s \ell^+ \ell^-$ processes. In the following discussion, we should assume that $\lambda'_{ij3} = \lambda_{ij2} = 0$.

Next, we will show the contributions of $R$-parity violating MSSM to $b \rightarrow s \ell^+ \ell^-$ processes. All the Feynman diagrams include four $W$ - $b$ box diagrams (Fig. 1a), five $W$ - $\tilde{b}_R$ box diagrams (one of which is Goldstone $-\tilde{b}_R$ box diagram) (Fig. 1b), one $H^\pm - \tilde{b}_R$ box diagram (Fig. 1c), two $\sqrt{2}$ box diagrams (Fig. 1d) and two $\gamma$-penguin diagrams (Fig. 2). Most of these results can be found in Refs [49, 50, 52, 58], however, to our knowledge, the results of the diagram induced by exchanging charged Higgs $H^\pm$ and right-handed sbottom $\tilde{b}_R$ in loop are the first to be given in this paper. The photon penguin diagrams, which have been neglected in previous work, play an important role in our discussion, as we will explain in more detail later. We do not find sizable $Z$-penguin contributions to $b \rightarrow s \ell^+ \ell^-$ processes. In this work, the contributions of $\gamma/Z$-penguin diagrams always include their supersymmetric counterparts unless otherwise specified. For convenience,

$$W_{\text{RPV}} = \mu_i L_i H_u + \frac{1}{2} \lambda'_{ijk} L_i L_j E^{c}_{k} + \lambda'_{ijk} L_i Q_j D^{c}_{k}$$

where the generation indices are denoted by $i, j, k = 1, 2, 3$ and the colour indices are suppressed. All repeated indices are assumed to be summed over throughout this paper unless otherwise stated (For example, repeated indices in both numerator and denominator are not automatically summed). $H_u, L$ and $Q$ are $SU(2)$ doublet chiral superfields while $E^c$, $D^c$ and $U^c$ are $SU(2)$ singlet chiral superfields.

In this work, we are mainly interested in the terms $\lambda'_{ijk} L_i Q_j D^{c}_{k}$ which related to both quarks and leptons. This choice can also alleviate the constraint of sneutrino masses on the collider, because the lower limit of sneutrino masses will be as high as TeV scale [62–65] when there are non-zero $\lambda$ and $\lambda'$ at the same time.

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the following Passarino-Veltman functions [67] $D_0$ and $D_2$
are defined as

$$D_0[m_1^2, m_2^2, m_3^2, m_4^2] = \int \frac{d^4k}{(2\pi)^4} \frac{1}{(k^2 - m_1^2)^2(k^2 - m_2^2)^2(k^2 - m_3^2)(k^2 - m_4^2)}$$

$$D_2[m_1^2, m_2^2, m_3^2, m_4^2] = \int \frac{d^4k}{(2\pi)^4} \frac{k^2}{(m_1^2 - m_2^2)(m_1^2 - m_3^2)(m_1^2 - m_4^2)}$$

The contributions of $W - \tilde{b}_R$ box diagram to $b \rightarrow s\mu^+\mu^-$ processes are given by

$$C_{9}^{V(W)} = \frac{-i\pi^2}{\sqrt{2}G_F \sin^2 \theta_W \eta_t} \times$$

$$\left( \lambda_{233}^\prime \lambda_{233}^{* V} V_{ub} D_2[m_W^2, m_{\tilde{b}_R}^2, m_{\tilde{u}^*_R}^2, m_{\tilde{t}_R}^2] \right)$$

where $\lambda_{ij}^\prime$ are the NP particle here. In the limit $m_{\tilde{b}_R} \gg m_t$, one has $C_{9}^{V(W)} = \frac{m_W^2}{16\pi\alpha m_{\tilde{t}_R}^2} |\lambda_{233}^\prime|^2$ [49, 50, 61] which is obviously positive.

The contributions of $H^\pm - \tilde{b}_R$ box diagram to $b \rightarrow s\mu^+\mu^-$ processes are given by

$$C_{9}^{V(H^\pm)} = \frac{-i\pi^2}{\sqrt{2}G_F \sin^2 \theta_W \tan^2 \beta \eta_t} \times$$

$$\lambda_{233}^{\lambda_{233}^\prime V} V_{ub} D_0[m_{H^\pm}^2, m_{\tilde{b}_R}^2, m_{\tilde{u}^*_R}^2, m_{\tilde{t}_R}^2]$$

which should be considered in the following numerical analysis. The tan $\beta = v_u/v_d$ where $v_u$ and $v_d$ are the vacuum expectation values of two Higgs doublets respectively.

The contributions of $4\lambda^\prime$ box diagram to $b \rightarrow s\mu^+\mu^-$ processes are given by

$$C_{9}^{V(4\lambda^\prime)} = \frac{-i\pi \lambda_{233}^{\lambda_{233}^\prime V}}{4\sqrt{2}G_F \alpha \eta_t} \times$$

$$\left( \lambda_{233}^{\lambda_{233}^\prime V} V_{ub} D_2[m_W^2, m_{\tilde{b}_R}^2, m_{\tilde{u}^*_R}^2, m_{\tilde{t}_R}^2] \right)$$

where $\lambda_{ij}^{\lambda_{ij}^\prime}$ are the NP particle here. In the limit $m_{\tilde{b}_R} \gg m_t$, one has $C_{9}^{V(4\lambda^\prime)}$ naturally gives a nonzero $C_{9}$.

The contributions of photonic penguin diagrams are leptonic flavour universal which naturally gives us a nonzero $C_{9}$.

$$C_{9}^{U} = \frac{\sqrt{2} \lambda_{233}^{\lambda_{233}^\prime V}}{36G_F \eta_t} \left[ \frac{1}{6m_{\tilde{b}_R}^2} - \left( \frac{4}{3} + \log \frac{m_W^2}{m_{\tilde{t}_R}^2} \right) \frac{1}{m_{\tilde{t}_R}^2} \right]$$
As stated in Ref. [52], this result is consistent with that in Ref. [68], but it is a negative sign different from that in Ref. [50]. The first term in Eq. (20) comes from the contribution of Fig. 2b, like the photonic penguin induced by scalar leptoquark. We find this term give a negligible contribution, which is in agreement with Refs. [61, 69]. However the second term in Eq. (20) has a significant contribution because of the logarithmic enhancement, which has never been addressed before. These photonic penguins also contribute new electromagnetic dipole operator $O_7 = \frac{2}{f^2 (8\pi^3) P_{\mu \nu}} F_{\mu \nu}$, which is strictly constrained by $B \to X_s \gamma$ decay [9]. Fortunately, we find that the corresponding contribution can be ignored numerically because there such logarithmic enhancement absent [50, 52, 68].

We will discuss the possibility of using muon sneutrinos $\tilde{\nu}_\mu$ and right-handed sbottoms $b_R$ to explain $b \to s\ell^+\ell^-$ anomaly, for which we set the mass of tauon sneutrinos $\tilde{\nu}_\tau$ and three left-handed up type squarks $u_{ij}$ sufficiently large that the contributions of the loop diagrams containing them are ignored. The contribution from $H^\pm - b_R$ box diagram is usually positive, but we can simply suppress this effect by increasing parameter $\tan \beta$. Thus, the contributions to only muon channel are

$$C^N_9 = -\frac{\sqrt{2} \lambda^2_{133} \lambda^2_{223} f(x_{\nu_\mu}) + |\lambda^2_{233}|^2 x_{b_R}}{32 G_F \sin^2 \theta_W m^2_{b_R}} \frac{16 \pi a}{|x_{b_R}|},$$

where $x_{\nu_\mu} \equiv m^2_{\nu_\mu}/m^2_W$, $x_{b_R} \equiv m^2_t/m^2_{b_R}$, and the loop function $f(x) \equiv x/(1-x+\log x)$.

### III. $R(D^{(*)})$ ANOMALY AND OTHER CONSTRAINTS

In this section, we discuss the interpretation of $R(D^{(*)})$ anomaly and consider the constraints imposed by other related processes from $B$, $D$, $K$, $\tau$, and $Z$ decays.

#### III.1. $R(D^{(*)})$ anomaly

In $R$-parity violating MSSM, the charged current processes $d_j \to u_n \ell_\nu \nu_i$ are induced by exchanging $b_R$ at tree level. The effective Lagrangian of these processes are given by

$$\mathcal{L}_{\text{eff}} = -\frac{4 G_F}{\sqrt{2}} V_{ij} (\delta_{ij} + C_{njli} \bar{u}_n \gamma_\mu P_L d_j \bar{l}_\nu \gamma_\mu P_L \nu_i) + \text{H.c.},$$

where the Wilson coefficient $C_{njli}$ is

$$C_{njli} = \frac{\lambda^2_{i323} \lambda_{jl3}}{4 \sqrt{2} G_F V_{nj} m^2_{b_R}}.$$

Because taking $\lambda^2_{i33} = 0$ to eliminate the contributions of box diagrams to $b \to s \ell^+\ell^-$ processes, we have $C_{njli} = C_{njli} = 0$. It is useful to define the ratio

$$R_{njl} \equiv \frac{B(d_j \to u_n l \nu)}{B(d_j \to u_n l \nu)_{\text{SM}}} = \sum_{i=1}^{3} |\delta_{li} + C_{njli}|^2,$$

and we have

$$\frac{R(D)}{R(D)_{\text{SM}}} = \frac{R(D^*)_{\text{SM}}}{R(D^*)} = \frac{2R_{233}}{R_{232} + 1}.$$ (25)

To obtain the allowed parameter region, we use the following best fit value in the $R$-parity violating scenario

$$\frac{R(D)}{R(D)_{\text{SM}}} = \frac{R(D^*)_{\text{SM}}}{R(D^*)} = 1.14 \pm 0.04.$$ (26)

#### III.2. Constraints from the tree-level processes

In the scenario we set up, some other processes receive tree level $R$-parity violating contributions. Here we mainly discuss the constraints from neutral current processes $B \to K^{(*)}\nu \bar{\nu}$, $B \to \pi \nu \bar{\nu}$, $K \to \pi \nu \bar{\nu}$, $D^0 \to \mu^+ \mu^-$ and $\tau \to \mu \nu \bar{\nu}$, as well as charged current processes $B \to 3\nu \bar{\nu}$, $D_s \to 3\nu \bar{\nu}$ and $\tau \to K \nu$. These decays relate to

$$\frac{\lambda^2_{i33} \lambda^2_{lm3} \bar{u}_m \gamma_\mu P_L d_j \bar{l}_\nu \gamma_\mu P_L \nu_i}{2 m^2_{b_R}},$$

$$\frac{\lambda^2_{i33} \lambda^2_{lm3} \bar{u}_m \gamma_\mu P_L d_j \bar{l}_\nu \gamma_\mu P_L \nu_i}{2 m^2_{b_R}},$$

$$\frac{\lambda^2_{i33} \lambda^2_{lm3} \bar{u}_m \gamma_\mu P_L d_j \bar{l}_\nu \gamma_\mu P_L \nu_i}{2 m^2_{b_R}}.$$ (29)

The effective Lagrangian for $B \to K^{(*)}\nu \bar{\nu}$, $B \to \pi \nu \bar{\nu}$ and $K \to \pi \nu \bar{\nu}$ decays are defined by

$$\mathcal{L}_{\text{eff}} = (C^\text{SM}_{mj} \delta_{ij} + C^\text{SM}_{mj} \delta_{ij} \bar{d}_m \gamma_\mu P_L d_j \bar{l}_\nu \gamma_\mu P_L \nu_i + \text{H.c.},$$

where

$$C^\text{SM}_{mj} = \frac{\lambda_{i33} \lambda_{lm3}}{2 m^2_{b_R}},$$ (31)

is the SM one. The loop function $X(x) \equiv x/(x + x^2 + 3x + 2)/(x - 1)$ with $x \equiv m^2_t/m^2_{b_R}$. The $R$-parity violating contributions are given by

$$C^\nu_{mj} \bar{\nu}_i = \frac{\lambda^2_{i33} \lambda^2_{lm3}}{2 m^2_{b_R}}.$$ (32)

$^2$ In fact, by combining the assumptions $\lambda^2_{i33} = 0$ and $\lambda^2_{ij1} = \lambda^2_{ij2} = 0$, we can get $\lambda^2_{i33} = 0$, which implies that the contribution of box diagrams of NP to the first generation leptons and sleptons is zero, because we only consider the terms $\lambda^2_{i33} \bar{l}_\nu Q J_D$. 

$^3$ For simplicity, we assume $\lambda^2_{i33} = 0$. 

$^4$ Note that $\lambda^2_{i33} \equiv 0$ and $\lambda^2_{ij1} \equiv \lambda^2_{ij2} \equiv 0$, we can get $\lambda^2_{i323} = 0$, which implies the contribution of box diagrams to $b \to s \ell^+\ell^-$ processes. 

$^5$ Since we only consider the terms $\lambda^2_{i33} \bar{l}_\nu Q J_D$, 

$^6$ The contributions of box diagrams to $b \to s \ell^+\ell^-$ processes.

$^7$ Since we only consider the terms $\lambda^2_{i33} \bar{l}_\nu Q J_D$, 

$^8$ The contributions of box diagrams to $b \to s \ell^+\ell^-$ processes.
It is useful to define the ratio

$$R_{m_j}^{\mu \nu} = \frac{B(d_j \to d_{m} \bar{\nu})}{B(d_j \to d_{m} \bar{\nu})_{\text{SM}}} \sum_{i=1}^{3} \left| C_{m_j}^{\mu \nu} \right|^2 + \sum_{i=1}^{3} \left| C_{m_j}^{\mu \nu} \right|^2 \right|_{\text{SM}}^2 .$$

(33)

The upper limit of $B \to K^{+} \pi^{-} \nu \bar{\nu}$ decay corresponds to $R_{m_j}^{\mu \nu} < 5.2$ [59, 71] at 95% confidence level (CL), and the upper limit of $B \to \pi^{0} \pi^{-} \nu \bar{\nu}$ decay is related to $R_{m_j}^{\mu \nu} < 830.5$ [72, 73] at 90% CL. By combining the SM prediction $B(K^+ \to \pi^+ \nu \bar{\nu})_{\text{SM}} = (9.24 \pm 0.83) \times 10^{-11}$ [74] with experimental measurement $B(K^+ \to \pi^+ \nu \bar{\nu})_{\text{exp}} = (1.7 \pm 1.1) \times 10^{-10}$ [75], we obtain a stringent constraint from $K \to \pi^{0} \nu \bar{\nu}$ decay that makes

$$|\lambda'_{23}^{\mu \nu}| < 7.4 \times 10^{-4} \left( m_{\text{br}}^2 / 1 \text{ TeV} \right)^2 .$$

(34)

Therefore, we will assume $\lambda'_{\nu i k} = 0$ to satisfy this constraint. At the same time, under this assumption, $B \to \pi \nu \bar{\nu}$ decay is unaffected by the NP.

The branching fraction for $D^0 \to \mu^+ \mu^-$ decay is given by [58]

$$B(D^0 \to \mu^+ \mu^-) = \frac{\tau_D f_D^2 m_D m_{\mu}^2}{32\pi} \left| \frac{\lambda'_{23}^{\mu \nu}}{2 m_{\text{br}}^2} \right|^2 \sqrt{1 - \frac{4 m_{\mu}^2}{m_D^2}} ,$$

(35)

where decay constant of $D^0$ is $f_D = 209.0 \pm 2.4$ MeV [76]. The mean life $\tau_D = 410.1 \pm 1.5$ fs [75] and the upper limit of branching fraction of $D^0 \to \mu^+ \mu^-$ decay is $6.2 \times 10^{-9}$ at 90% CL [75]. The corresponding constraint is $|\lambda'_{23}^{\mu \nu}| < 0.31 \left( m_{\text{br}}^2 / 1 \text{ TeV} \right)^2$.

The branching fraction for $\tau \to \mu \rho^0$ decay is given by [77]

$$B(\tau \to \mu \rho^0) = \frac{\tau \rho f_{\rho}^2 m_{\rho}^3}{128\pi} \left| \frac{\lambda'_{13}^{\mu \nu}}{2 m_{\text{br}}^2} \right|^2 \left( 1 - \frac{m_{\rho}^2}{m_\tau^2} \right) \left( 1 + \frac{m_{\rho}^2}{m_\tau^2} - 2 \frac{m_{\rho}^2}{m_\tau^2} \right) ,$$

(36)

where $\tau = 290.3 \pm 0.5$ fs and the decay constant $f_{\rho} = 153$ MeV [50]. The current experimental upper limit on the branching fraction for this process is $B(\tau \to \mu \rho^0) < 1.2 \times 10^{-8}$ at 90% CL [75]. The corresponding constraint is $|\lambda'_{23}^{\mu \nu}| < 0.38 \left( m_{\text{br}}^2 / 1 \text{ TeV} \right)^2$.

The formulas for charged current processes are given, respectively, by

$$\frac{B(B \to \tau \nu)}{B(B \to \tau \nu)_{\text{SM}}} = R_{133} ,$$

(37)

$$\frac{B(D_s \to \tau \nu)}{B(D_s \to \tau \nu)_{\text{SM}}} = R_{223} ,$$

(38)

$$\frac{B(\tau \to K \nu)}{B(\tau \to K \nu)_{\text{SM}}} = R_{123} .$$

(39)

The corresponding experimental and theoretical values are listed, respectively, as follows: $B(B \to \tau \nu)_{\text{exp}} = (1.09 \pm 0.24) \times 10^{-4}$ [75], $B(B \to \tau \nu)_{\text{SM}} = (9.47 \pm 1.82) \times 10^{-5}$ [78]; $B(D_s \to \tau \nu)_{\text{exp}} = (5.48 \pm 0.23)%$ [75], $B(D_s \to \tau \nu)_{\text{SM}} = (5.40 \pm 0.30)%$; $B(\tau \to K \nu)_{\text{exp}} = (6.96 \pm 0.10) \times 10^{-3}$ [75], $B(\tau \to K \nu)_{\text{SM}} = (7.15 \pm 0.026) \times 10^{-3}$ [56].

III.3. Constraints from the loop-level processes

First of all, the most important one-loop constraint comes from $B_s - \bar{B}_s$ mixing, which is governed by

$$\mathcal{L}_{\text{eff}} = (C_{\text{SM}} + C_{\text{NP}})(\bar{\nu}_\mu P_L b)(\bar{\nu}_\mu P_L b) + \text{h.c.} ,$$

(40)

where the SM and NP Wilson coefficients are given respectively by

$$C_{\text{SM}} = -\frac{1}{4\pi^2} G_F^2 m_W^2 \epsilon_i^2 S(x_i) ,$$

(41)

$$C_{\text{NP}} = -\frac{1}{128\pi^2} \left[ (\lambda'_{13}^{\mu \nu} \lambda'_{23}^{\mu \nu})^2 + (\lambda'_{23}^{\mu \nu} \lambda'_{23}^{\mu \nu})^2 \right] ,$$

(42)

where loop function $S(x_i) = x_i (4 - 11 x_i + x_i^2) + 3 x_i^2 \log(3 x_i)$. At 2$\sigma$ level, the UTfit collaboration [79] gives the bound $0.93 < |1 + C_{\text{NP}} / C_{\text{SM}}| < 1.29$.

Next, we investigate a series of $Z$ decaying to two charged leptons with the same flavour like $Z \to \mu \mu (\tau \tau)$ and the different one like $Z \to \mu \tau$. The amplitude of these diagrams is $i M = i \frac{4 \times 11 x_i + x_i^2}{(4 x_i + 1)^3} B_{ij}^3 \epsilon_i^2 \epsilon_j^3 \bar{\nu}_\mu P_L \nu_\tau \epsilon_j$ [50], where $B_{ij}^3 = B_{ij}^1 + B_{ij}^2$ and

$$B_{ij}^1 = \sum_{l=1}^{2} \lambda'_{1l} \lambda'_{33} \frac{m_Z^2}{m_l^2} \left( 1 - \frac{4}{3} \sin^2 \theta_W \right) \times \left( \log \frac{m_Z^2}{m_l^2} - \frac{1}{3} \right) + \frac{\sin^2 \theta_W}{9} \right) ,$$

(43)

$$B_{ij}^2 = \frac{3 \lambda'_{13} \lambda'_{33}}{16 m_Z^2} \left( 1 + \log x_{br} \right) \times \left( 11 - 10 \sin^2 \theta_W \right) + (6 - 8 \sin^2 \theta_W) \log x_{br}$$

$$+ \frac{1}{10} \left( -9 + 16 \sin^2 \theta_W \right) \frac{m_Z^2}{m_l^2} \right) ,$$

(44)

here $B_{ij}^1$ is the contribution from the diagram induced by exchanging $b_R - u - u$ or $b_R - c - c$ in triangular loop and $B_{ij}^2$ is the contribution from the diagram induced by exchanging $b_R - t - t$ in triangular loop. As shown in Ref. [50], for $Z \to \mu \mu (\tau \tau)$, demanding the interference term in the partial width between the SM tree-level contribution and the NP one-loop level ones is less than twice the experimental uncertainty on the partial width [75], there are the bounds $|R(B_{22})| < 0.32$ and $|R(B_{33})| < 0.39$ [50]. And the experimental upper limit $B(Z \to \mu \nu) < 1.2 \times 10^{-5}$ [75] makes the bound $\sqrt{B_{22}^0 + B_{33}^0} < 2.1$ [50].

Finally, we discuss the lepton-flavour violating decay of $\tau$ lepton, including $\tau \to \mu \nu$ and $\tau \to \mu \mu$. In the limit
\[ B(\tau \rightarrow \mu \gamma) = \frac{\tau \alpha m^5}{4} (|A_2^L|^2 + |A_2^R|^2), \]

where the effective couplings \( A_{2}^{L,R} \) come from on shell photonic penguin diagrams [68].

\[ A_2^L = -\frac{\lambda'_{233} \lambda_{333}^{\ast}}{64 \pi \tau m_b^2}, \quad A_2^R = 0. \]

The current experimental upper limit is \( B(\tau \rightarrow \mu \gamma) < 4.4 \times 10^{-8} \) at 90% CL [75].

In general, the effective Lagrangian leading to \( \tau \rightarrow \mu \mu \) decay is given by [80, 81]

\[ \mathcal{L}_{\text{eff}} = -\frac{B_1}{2} (\bar{\tau} \gamma^\nu P_L \mu)(\bar{\mu} \gamma^\nu P_R \mu) - \frac{B_2}{2} (\bar{\tau} \gamma^\nu P_R \mu)(\bar{\mu} \gamma^\nu P_L \mu) + C_1(\bar{\tau} P_R \mu)(\bar{\mu} P_L \mu) + C_2(\bar{\tau} P_L \mu)(\bar{\mu} P_R \mu) + G_1(\bar{\tau} \gamma^\nu P_R)(\bar{\mu} \gamma^\nu P_L) + G_2(\bar{\tau} \gamma^\nu P_L)(\bar{\mu} \gamma^\nu P_R) - A_R (\bar{\tau} [\gamma, \gamma, \mu](\bar{\mu} \gamma^\nu P_R) + A_L (\bar{\tau} [\gamma, \gamma, \mu](\bar{\gamma} \mu P_L)) + H.c. \]

This Lagrangian leads to [80, 81]

\[ B(\tau \rightarrow 3\mu) = \frac{\tau \alpha m^5}{64 \pi \tau \rho^3} \left[ |B_1|^2 + |B_2|^2 + 8(|G_1|^2 + |G_2|^2) + \frac{|C_1|^2 + |C_2|^2}{2} + 32 \left(4 \log \frac{m^2}{m^2} - 11\right) \frac{|A_R|^2 + |A_L|^2}{m^2} - \frac{32}{m^2} (A_L G_1^2 + A_R G_2^2) \right]. \]

In our scenario, there are three different types of contributions, the photonic and Z penguins as well as box diagrams with four \( \chi \) couplings, that can contribute to \( \tau \rightarrow \mu \mu \) decay. The nonzero Wilson coefficients are [50, 68]

\[ B_1 = -2(4 \pi \alpha A_2^L + \sin^2 \theta_W B') \],

\[ B_2 = 4 \pi \alpha A_1^L + \left( -\frac{1}{2} + \sin^2 \theta_W \right) B' + C_\tau, \]

\[ A_L = 2 \pi \alpha m_\tau A_2^L \]

where

\[ B' = -\frac{3 \lambda'_{233} \lambda_{333}^{\ast} x_R}{8 \pi \cos \theta_W \sin^2 \theta_W m^2}, \]

\[ C_\tau = \sqrt{\frac{\lambda'_{233} \lambda_{333}^{\ast} \lambda_{23}^{\ast}}{4}} \frac{1}{4 \pi \tau m^2} \left[ \left( 1 + \frac{1}{2} \frac{1}{3} \frac{4}{3} \log \frac{m^2}{m^2} \right) \right], \]

and the off-shell effective coupling \( A_{2}^{L} \) is [68]

\[ A_2^L = \frac{\lambda'_{233} \lambda_{333}^{\ast} \lambda_{23}^{\ast}}{16 \pi \tau m^2} \left[ \frac{1}{18} - \frac{2}{3} \frac{1}{3} + \frac{4}{3} \log \frac{m^2}{m^2} \right]. \]

The current experimental upper limit on the branching fraction for this decay is \( B(\tau \rightarrow \mu \mu) < 2.1 \times 10^{-8} \) at 90% CL [75].

### IV. NUMERICAL RESULTS AND DISCUSSIONS

In this section, we discuss how to interpret both \( b \rightarrow s l^+ l^- \) and \( R(D^{(*)}) \) anomalies and satisfy all these potential constraints simultaneously. The relevant model parameters in our scenario are the wino mass \( m_{\tilde{W}} \), the mass of muon sneutrino \( m_{\tilde{\mu}} \), and the mass of right-handed sbottom \( m_b^R \), as well as four nonzero couplings \( \lambda'_{233}, \lambda'_{333}, \lambda'_{23}, \lambda_{333} \). We set \( m_\tilde{W} = 270 \text{ GeV} \). It can be seen from Ref. [52] that a positive product \( \lambda'_{233} \lambda_{333} \) is needed to explain the \( b \rightarrow s l^+ l^- \) anomaly mainly through muon sneutrinos (the \( C^\nu_{\mu} \) part). Both \( \lambda'_{233} \) and \( \lambda_{333} \) are positive to help solve \( R(D^{(*)}) \) anomaly by exchanging \( b_R \) at tree level [56]. The combination of the choice of above couplings will naturally produce a negative \( C^\nu_{\mu} \), which is in line with the conclusion of the global analysis [20]. Our numerical results are shown in Fig. 3. These results show that it is possible to explain \( b \rightarrow s l^+ l^- \) and \( R(D^{(*)}) \) anomalies simultaneously at 90% CL. The regions of NP parameters that can solve B-physics anomalies are most constrained by \( B \rightarrow K^{(*)} \nu \bar{\nu} \) decays and \( B_s - \bar{B}_s \) mixing. In addition, the processes of Z decays can provide a weak constraints. We find that other related processes, such as \( D^0 \rightarrow \mu^+ \mu^- \), \( \tau \rightarrow \mu \), \( B \rightarrow \tau \nu \), \( D_s \rightarrow \tau \nu \), \( \tau \rightarrow K \nu \), \( \tau \rightarrow \mu \gamma \), and \( \tau \rightarrow \mu \mu \) decays, do not provide available constraints.

We show in Fig. 3a and Fig. 3b the allowed regions in the planes of coupling parameters \( (\lambda'_{233}, \lambda_{333}) \) and \( (\lambda'_{23}, \lambda_{333}) \) respectively when other parameters are fixed. These two subfigures show that in order to explain the B-physics anomalies, the coupling parameters need to satisfy the relation \( \lambda_{333} > \lambda'_{233} > \lambda'_{23} \approx \lambda_{333} \), and the required \( \lambda'_{233} \) and \( \lambda_{333} \) are very small. Therefore, the next four subfigures in Fig. 3 mainly discuss the relationships between the coupling parameters \( \lambda'_{233} \) and \( \lambda_{333} \) and the masses \( m_{\tilde{W}} \) and \( m_{\tilde{\mu}} \). From Fig. 3a, we can see that \( \lambda'_{233} \) is more constrained by \( R(D^{(*)}) \), \( B \rightarrow K^{(*)} \nu \bar{\nu} \) decay and Z decays, but less affected by \( b \rightarrow s l^+ l^- \) processes and \( B_s - \bar{B}_s \) mixing. On the contrary, \( \lambda_{333} \) is greatly constrained by \( b \rightarrow s l^+ l^- \) processes and \( B_s - \bar{B}_s \) mixing, but has little influence on \( R(D^{(*)}) \), \( B \rightarrow K^{(*)} \nu \bar{\nu} \) decay and Z decays. As shown in Fig. 3c, after the variable parameter \( m_{\tilde{W}} \) is added, the constraints of \( \lambda'_{233} \) from \( R(D^{(*)}) \), \( B \rightarrow K^{(*)} \nu \bar{\nu} \) decay and Z decays will be relaxed a lot. The parameters \( \lambda_{333} \) and \( m_{\tilde{W}} \) are highly correlated. Because we choose a smaller mass of muon sneutrino, the \( B_s - \bar{B}_s \) mixing is more sensitive to \( m_{\tilde{\mu}} \) than to \( m_{\tilde{W}} \), which can be seen by comparing Fig. 3c with Fig. 3e, or Fig. 3d with Fig. 3f. All subfigures contain parameter spaces (marked in purple) that can resolve \( b \rightarrow s l^+ l^- \) and \( R(D^{(*)}) \) anomalies, and satisfy the constraints from other related processes simultaneously.

### V. CONCLUSIONS

The recent measurements on semileptonic decays of B-meson suggest the existence of NP which breaks the LFU. Among them, the observables \( R_K^{\nu} \) and \( P_{f}^{\nu} \) in \( b \rightarrow s l^+ l^- \) processes and the \( R(D^{(*)}) \) in \( B \rightarrow D^{(*)}\nu \nu \) decays are
more striking. They are collectively called B-physics anomalies. In this work, we have explored the possibility of using muon sneutrinos \( \tilde{\nu}_\mu \) and right-handed sbottoms \( b_R \) to solve these B-physics anomalies simultaneously in R-parity violating MSSM.

To explain the anomalies in \( b \to s \ell^+ \ell^- \) processes, we use a two-parameter scenario, where the total Wilson coefficients of NP are divided into two parts, one is the \( C_{9}^{U} \) (Noting \( C_{10}^{NP} = -C_{9}^{U} \)) that only contributes the muon channel and the other is the \( C_{9}^{U} \) that contributes both the electron and the muon channels. First, we scrutinize all the one-loop contributions of the superpotential terms \( \lambda_{ij}\bar{Q}_i Q_j D_{\ell}\) to the \( b \to s \ell^+ \ell^- \) processes under the assumptions \( \lambda'_{ij} = \lambda''_{ij} = 0 \) and \( \lambda_{ij} = 0 \). We find that the contribution from the \( H^+ - \bar{b}_R \) box diagram (Fig. 1c) is missed in the literature, this contribution is usually positive, and we can suppress it by increasing parameter \( \tan \beta \). The photonic penguin induced by exchanging sneutrino can provide important contribution due to the existence of logarithmic enhancement, which has never been addressed before. This contribution is lepton flavour universal due to the SM photon, so it is natural to contribute a nonzero \( C_{9}^{U} \).

Global analyses show that the sizable magnitude of \( C_{9}^{U} \) needed to explain \( b \to s \ell^+ \ell^- \) anomaly. However, \( C_{9}^{U} \) in the scenario with nonzero \( C_{9}^{U} \) is smaller than the one in the scenario without \( C_{9}^{U} \). With the addition of the latest measurements from the Belle collaboration, the world averages of \( R(D^{(*)}) \) are closer to the predicted values of the SM. These changes make it possible to use \( \tilde{\nu}_\mu \) and \( b_R \) to explain \( b \to s \ell^+ \ell^- \) and \( R(D^{(*)}) \) anomalies, simultaneously. We also consider the constraints of other related processes in our scenario. The strongest constraints come from \( B \to K^{(*)}\nu\bar{\nu} \) decays and \( B_s - \bar{B}_s \) mixing. Besides, the processes of \( Z \) decays can provide a few constraints. The other decays, such as \( D^{0} \to \mu^+\mu^- \), \( \tau \to \mu\rho^0 \), \( B \to \tau\nu \), \( D_s \to \tau\nu \), \( \tau \to K\nu \), \( \tau \to \mu\gamma \), and \( \tau \to \mu\mu\mu \), do not provide available constraints.

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