Probing \((g - 2)_\mu\) at the LHC in the paradigm of R-parity violating MSSM

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The measurement of the anomalous magnetic moment of the muon exhibits a long standing discrepancy compared to the Standard model prediction. In this paper, we concentrate on this issue in the framework of R-parity violating Minimal supersymmetric standard model. Such a scenario provides substantial contribution to the anomalous magnetic moment of the muon while satisfying constraints from low energy experimental observables as well as neutrino mass. In addition, we point out that the implication of such operators satisfying muon \((g - 2)\) are immense from the perspective of the LHC experiment, leading to a spectacular four muon final state. We propose an analysis in this particular channel which might help to settle the debate of R-parity violation as a probable explanation for \((g - 2)_\mu\).

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I. INTRODUCTION

We are living in an era enriched with many experimental breakthroughs and results. Recently, the two CERN based experiments, namely ATLAS and CMS collaborations have confirmed the existence of a neutral boson, widely accepted to be the Higgs boson with mass close to 125 GeV \([1]\). All the decay modes of this scalar boson have been measured with moderate accuracy and the results obtained so far are fairly consistent with the standard model (SM) expectation. However, from an aesthetic point of view, the SM inevitably has the hierarchy problem which is associated to the stabilization of the Higgs boson mass from large radiative corrections. Further, the observation of neutrino mass and mixing and the existence of dark matter (DM) most certainly require beyond the standard model (BSM) physics. Another sector which requires the intervention of BSM theories is the anomalous magnetic moment of the muon, quantified as \(a_\mu = (g - 2)_\mu/2\), which has been measured with unprecedented accuracy at the Brookhaven \((g - 2)\) experiment. However, there still exists a discrepancy between the experimental observation and the SM prediction, given by \(\Delta a_\mu = a_\mu^{\text{exp}} - a_\mu^{\text{SM}} = (29.3 \pm 9.0) \times 10^{-10}\) \([2]\). This anomaly with respect to the SM expectation reflects the contributions arising from the perturbative higher order electroweak corrections, the virtual hadronic inputs and the possible presence of the BSM physics.

Supersymmetry (SUSY) \([3–6]\) remains one of the most celebrated BSM theories till today. The minimal supersymmetric standard model (MSSM) provides an elegant solution to the hierarchy problem \([4, 6]\). In addition neutrino masses and DM can also be explained in the paradigm of MSSM. Another important feature of MSSM is that it yields sizeable contribution to the muon \((g - 2)\) requiring light first two generation of sleptons \([7, 8]\). However, the ATLAS and CMS experiments, in their hunt for superpartners, have found no significant excess over the SM background after the 7+8 TeV run of the LHC \([9, 10]\). For comparable gluino and first two generation squark masses, the bound on these particles can be as large as 1.7 TeV in R-parity conserving (RPC) and simplified phenomenological MSSM (pMSSM) scenario \([11]\). On the other hand, the constraints on the first two generation sleptons are comparatively weaker and lies in the ballpark of 300 GeV \([12, 13]\).

In MSSM, the loop contributions to \((g - 2)_\mu\) arises if there is a chirality transition in the external muon lines. This chirality transition requires an insertion of a fermion mass or a Yukawa coupling vertex. In the framework of R-parity conserving SUSY, the main possibilities for the chirality flip are the following, a) a muon line through a muon mass term, which contributes to a factor \(m_\mu\), b) a Yukawa coupling in between the Higgs field and \(\mu_L\) and \(\mu_R\), which contributes to a factor \(y_\mu\), c) a \(L – R\) mixing in the scalar sector, more precisely corresponding to a transition between \(\tilde{\mu}_L\tilde{\mu}_R\), which contributes to a factor proportional to \(m_{\tilde{\mu}}\tan\beta\), where \(\mu\) is the Higgsino mass parameter and \(\tan\beta\) is the ratio of two vacuum expectation values (vevs) \(v_u\) and \(v_d\) associated with the two Higgs doublets \(H_u\) and \(H_d\) respectively. Finally d) a SUSY Yukawa coupling of a Higgsino to muon and \(\tilde{\mu}\) or \(\tilde{\nu}_\mu\), contributing a factor of \(y_\mu\). It is evident that all of these contribute to the muon \((g - 2)\) and an overall rough estimate implies \(a_\mu \sim m^2_\mu/M^2_{\text{SUSY}}\) \([14]\). Therefore, the new physics scale or more precisely the SUSY scale must be around \(O(100)\) GeV, i.e., the electroweak scale to have large contributions to \((g - 2)_\mu\).

On the other hand, one of the many interesting outcomes of R-parity violating (RPV) MSSM \([15]\) is that it is an intrinsic way by which substantial augmentation

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of muon \((g - 2)\) can be obtained [16]. Further RPV is also interesting as it is an inherent SUSY way to generate neutrino masses both at the tree level as well as at the one loop level. In this work, we consider a RPV MSSM scenario with relevant operators, which can give sizeable contribution to the anomalous magnetic moment of the muons. We respect the collider bounds on the slepton masses as well as indirect constraints from neutrino masses and low energy observables to present a self consistent picture. Most importantly, RPV MSSM also provide direct spectacular signals at the LHC. It is important to note that we ignore the contributions originating from left scalar and right scalar fermion mixing terms as they are negligible. There exists several phenomenological studies incorporating the muon \((g - 2)\) anomaly and the LHC bounds in R-parity conserving and violating SUSY framework, a partial list can be seen in Refs. [17–26].

The plan of this paper is as follows. In section II, we look into the theoretical framework of the study under consideration and its effects on \((g - 2)_\mu\). In section III, we study the relevant constraints coming from low energy observables and neutrino masses, which is necessary for considering a \(\mathcal{O}(1)\) RPV coupling. Section IV is dedicated to a numerical analysis with appropriate benchmark points followed by a detailed discussion on the present bounds from LHC data. In section V, we perform a dedicated collider analysis to correlate the fact that \((g - 2)_\mu\) from the RPV MSSM scenario can leave its fingerprints in the LHC experiments. Concluding remarks and related discussions are relegated to section VI.

II. MUON \((g - 2)\) IN MSSM

When \(R\)-parity is conserved, the SUSY effects on \(a_\mu\) includes contribution from the chargino-muon sneutrino and neutralino-smuon loops. The generic expressions for one-loop SUSY contributions to \(a_\mu\), including the effects of possible complex phases are given as [27, 28]

\[
a_\mu^{\tilde{\chi}_0^0} = \frac{m_\mu}{16\pi^2} \sum_{i,m} \left\{ - \frac{m_\mu}{12m_{\tilde{\chi}_m^i}^2} (|n_{im}^L|^2 + |n_{im}^R|^2) F_1^N(x_{im}) + \frac{m_{\tilde{\chi}_m^i}^0}{3m_{\tilde{\chi}_m^i}^2} \text{Re}(n_{im}^L n_{im}^R) F_2^N(x_{im}) \right\},
\]

\[
a_\mu^{\tilde{\chi}_\pm^\pm} = \frac{m_\mu}{16\pi^2} \sum_k \left\{ \frac{m_\mu}{12m_{\tilde{\chi}_k^k}^2} (|c_k^L|^2 + |c_k^R|^2) F_1^C(x_k) + \frac{2m_{\tilde{\chi}_k^k}}{3m_{\tilde{\chi}_k^k}^2} \text{Re}(c_k^L c_k^R) F_2^C(x_k) \right\}
\]

where \(i, m = 1, 2, 3, 4\), \(m = 1, 2\) and \(k = 1, 2\) denotes the neutralino, smuon and chargino mass eigenstates respectively. The couplings are defined as

\[
n_{im}^R = \sqrt{2} g_1 N_{i1} X_{m2} + y_\mu N_{i3} X_{m1},
\]

\[
n_{im}^L = \frac{1}{\sqrt{2}} (g_2 N_{i2} + g_1 N_{i1}) X_{m1} - y_\mu N_{i3} X_{m2},
\]

\[
c_k^R = y_\mu U_{k2},
\]

\[
c_k^L = -g_2 V_{k1},
\]

where \(N\) represents neutralino, \(U\) and \(V\) are chargino mixing matrices respectively while \(X\) denotes the slepton mixing matrix. The muon yukawa coupling \(y_\mu = g_2 m_\mu/\sqrt{2} m_W \cos \beta\) and the kinematic loop functions are defined in terms of the variables \(x_{im} = m_{\tilde{\chi}_m^i}^2/m_{\tilde{\chi}_m^i}^2\) and

\[
x_k = m_{\tilde{\chi}_k^k}^2/m_{\tilde{\chi}_k^k}^2\] and are as follows

\[
F_1^N(x) = \frac{2}{(1 - x)^4} \left[ 1 - 6x + 3x^2 + 2x^3 - 6x^2 \ln x \right],
\]

\[
F_1^C(x) = \frac{3}{(1 - x)^3} \left[ 1 - x^2 + 2x \ln x \right],
\]

\[
F_2^N(x) = \frac{2}{(1 - x)^4} \left[ 2 + 3x - 6x^2 + x^3 + 6x \ln x \right],
\]

\[
F_2^C(x) = -\frac{3}{2(1 - x)^3} \left[ 3 - 4x + x^2 + 2 \ln x \right] .
\]

In the limit when all the mass scales are roughly of the same order, i.e., \(M_{\text{SUSY}}\), the sum of the above expressions in Eq. 1 and Eq. 2 reduces to a more simpler form as [27, 28] (see Appendix for more detail)

\[
[a_\mu^{\text{SUSY}}]_{\text{RPC}} \simeq 14 \tan \beta \left( \frac{100 \text{GeV}}{M_{\text{SUSY}}} \right)^2 10^{-10} .
\]

Furthermore, as we have already discussed, in the absence of \(R\)-parity, the superpotential contains additional terms which are lepton and baryon number violating. In the context of our analysis we will consider only the fol-
lowing terms in the superpotential

\[ W_{\mathcal{L}_\mu} = W_{\text{MSSM}} + \frac{1}{2} \lambda_{ijk} \tilde{l}_i \tilde{l}_j \tilde{E}_k, \]  

(6)

where \( W_{\text{MSSM}} \) contains the usual MSSM superfields and \( \tilde{l}, \tilde{E} \) are the left-chiral lepton and left-chiral anti-lepton superfields respectively. Gauge invariance enforces \( \lambda_{ijk} \) to be antisymmetric with respect to their first two indices. As a result, \( \lambda_{ijk} = -\lambda_{jik} \). These RPV terms in the superpotential yield the following terms in the Lagrangian in the form as

\[ \mathcal{L} = -\frac{1}{2} \lambda_{ijk} \left[ \tilde{\nu}_L \tilde{l}_i R \tilde{l}_j L + \tilde{\nu}_L \tilde{l}_i j \nu L + \tilde{l}_i k R \nu \tilde{\nu}_R j L - (i \leftrightarrow j) \right] + \text{h.c.} \]  

(7)

In the four component notation, the terms in Eq. 7 which contribute to \((g - 2)_{\mu}\) can be explicitly written as

\[ \mathcal{L} \subset -\lambda_{ij2} \left[ \tilde{\nu}_L \tilde{l}_i i P l_j + \tilde{\nu}_L \tilde{l}_i j P l_i \right] - \lambda_{i2k} \left[ \tilde{\nu}_L \tilde{l}_i k P l_k + \tilde{l}_i k R P l_k \right] + \text{h.c.} \]  

(8)

In our scenario, we assume the first two generations of right chiral sleptons to be heavy to evade constraints appearing from neutrino masses and low energy observables in the presence of order one \( \lambda' \)’s as elaborated later. However, the left chiral charged sleptons/sneutrinos can be light and still avoiding bounds from direct collider constraints. Therefore, in addition to the \( \Delta a_{\mu} \) contribution coming from RPV operators, we also have sizeable contribution from the RPC section. The relevant diagrams contributing to \((g - 2)_{\mu}\) are shown in Fig. 1. The generic expression of \((g - 2)_{\mu}\) in the context of RPV MSSM are written as [16]

\[ \left[ a_{\mu}^\lambda \right]_{\text{RPV}} = \frac{m_{\mu}^2}{96\pi^2} \left[ \lambda_{23k}^2 \frac{2}{m_{\nu}^2} + \left| \lambda_{3k2} \right|^2 \frac{2}{m_{\nu}^2 - \frac{1}{m_{\tau}^2}} \right] \]  

(9)

\[ - \left| \lambda_{k23} \right|^2 \frac{1}{m_{\tau}^2}, \]

where \( m_{\tilde{\tau}} \) and \( m_{\tilde{\nu}} \) are the left and right chiral stau masses respectively while \( m_{\nu} \) is the tau-sneutrino mass. In the limit where all the relevant third generation slepton masses are considered to be equal, i.e., \( m_{\tilde{\tau}} = m_{\tilde{\nu}} = m_{\tilde{\nu}} = \tilde{m} \), then Eq. (9) reduces to the following simplified form

\[ \left[ a_{\mu}^\lambda \right]_{\text{RPV}} = \frac{m_{\mu}^2}{32\pi^2 m_{\tilde{\nu}}^2} \left[ \frac{1}{3} \left| \lambda_{312} \right|^2 + \frac{2}{3} \left| \lambda_{321} \right|^2 + \frac{1}{3} \left| \lambda_{323} \right|^2 \right] \]  

(10)

An important observation is except \( \lambda_{322} \), all the other RPV couplings come with a factor less than one. Our goal is now to study the present bounds on these couplings and to make sure if such an order one \( \lambda \) can be considered. We note in passing that in the present work we have taken into account all the contributions to anomalous magnetic moment of the muon coming from both the RPC as well as RPV MSSM.

### III. Bounds on RPV Couplings

In the paradigm of SM, the lepton flavor violating (LFV) processes occur at a negligible rate due to the smallness of the neutrino masses. As a result, they are sensitive probe of new physics and can be used to place
bounds on $R'_p$ couplings\(^2\). In order to disentangle the effects of $R'_p$ interactions from the effects emerging from the possible flavor non-universalities in the scalar lepton sector, we assume that the slepton mass matrices are diagonal with first two generations having equal masses. The possible sources of LFV are noted down in the following processes.

- Lepton flavour violating radiative decays of charged leptons : The $R'_p$ interactions can in principle generate LFV decays of charged leptons, such as $l_i \to l_j \gamma$ through one loop diagrams\([45]\).
- Lepton flavour violating decays of $\mu$ and $\tau$ into three charged leptons : The LFV decays like $l_m \to l_i l_j l_k$, where $l_m = \mu, \tau$ can be mediated at the tree level through $t$ and $u$-channel sneutrino exchanges when the involved leptons posses non-zero $\lambda$ type $R'_p$ couplings. The non-observation of these processes results in bounds on $\lambda_{\alpha m i} \lambda_{\alpha j k}$, where the sneutrino carries the index $n$\([46]\).
- Muon to electron conversion in nuclei : $\mu^- \to e^- \nu_e$ conversion in a nucleus is normally induced by $\lambda \lambda'$ or $\lambda' \lambda'$ couplings\(^3\). However, $\mu^- \to e^- \nu_e$ conversion in a nucleus can also proceed through photon penguin diagrams. The associated bounds can be much stronger than the ones extracted from the previously mentioned processes. The non-observation of these processes can be translated into bounds on $\lambda \lambda'$ couplings\([48–50]\).
- Charged current universality and bounds from $R_e/R_{\mu}$ : One should also take into account bounds from charged current universality which results in single bounds on the $\lambda$ couplings. Similar bounds can also be obtained from the ratio $R_e = \Gamma(\tau \to e\nu\nu)/\Gamma(\tau \to \mu\nu\nu)$ and $R_{\mu} = \Gamma(\tau \to \mu\nu\nu)/\Gamma(\mu \to e\nu\nu)$\([45]\).

We now tabulate the bounds on the relevant $R'_p$ couplings from the non-observation of the processes as mentioned earlier. All these limits are obtained from BP1 of $R'_p$ couplings from low energy experiments\([45, 46, 48–50]\) with specific benchmark point as shown in Ref.\([29]\).

| $R'_p$ couplings | $l_i \to l_j \gamma$ | $l_i \to 3l_j$ | $\tau \to l_i P/\mu - e$ | $l_i \to l_j l_k l_l$ |
|------------------|---------------------|----------------|---------------------|---------------------|
| $|\lambda_{312}^\alpha l_{322}|$ | $2.3 \times 10^{-3}$ | $8.2 \times 10^{-4}$ | $1.6 \times 10^{-4}$ | - |
| $|\lambda_{321} l_{322}|$ | $3.8 \times 10^{-4}$ | $4.1 \times 10^{-4}$ | $1.2 \times 10^{-4}$ | - |
| $|\lambda_{323} l_{322}|$ | - | - | - | $2.4 \times 10^{-4}$ |

\(\text{TABLE I: Bounds on } R'_p \text{ couplings from low energy experiments [45, 46, 48–50] with specific benchmark point as shown in Ref.\([29]\).}

\(^2\) For bounds on trilinear R-parity violating couplings see Refs.\([29–44]\).

\(^3\) For a theoretical calculation of this process in R-parity-violating SUSY models, we refer Ref.\([47]\).

Ref\([29]\), where the first two generations are considered to be heavy with masses around 1 TeV and the third generation is light. Making the first two generations masses heavier would further relax the bounds on $R'_p$ couplings. However, we take a more conservative approach here and use the strongest limits. In addition, from the charge current universality one finds $|\lambda_{123}| \sim 0.049 \times m_{\tilde{\mu}}/100 \text{GeV}$.

Since, in our framework, the third generation is considered to be light hence the bounds on the particular $R'_p$ operators turn out to be stringent and should be respected. From $R_e$ and $R_{\mu}$, one finds $|\lambda_{322}| < 0.07 \times m_{\tilde{\mu}}/100 \text{ GeV}$. Hence, in our scenario this bound can be readily relaxed by assuming large mass for the second generation sleptons which we have considered. As mentioned earlier, this bound also translates to a lower bound on $\tilde{e}_R$ as they are considered to be degenerate with $\tilde{\mu}_R$.

From the above discussion and Table I, it is conspicuous that only one of the $R'_p$ violating operator can be large ($O(1)$) satisfying the above mentioned constraints. We choose it to be $\lambda_{322}$.

- Bound on $R'_p$ couplings from neutrino mass : Neutrino masses provide serious constraints on trilinear $R'_p$ couplings. In this section we will compute the impact of neutrino masses on $\lambda_{322}$. These couplings generate neutrino masses radiatively (see Fig. 2) and the generic expression is noted down as\([51, 52]\)

$$m_{\nu} \sim \frac{1}{32\pi^2} \sum_{i,p} \lambda_{qlp} \lambda_{mpq} m_i \sin 2\phi_l \ln \left(\frac{M_{\mu}^2}{M_{\nu}^2}\right)$$

where $m_i$ is the mass of the lepton, $\phi_l$ is the mixing angle obtained by diagonalising the charged slepton mass squared matrix, which takes the form

$$\sin 2\phi_l = \frac{2A m_i}{\sqrt{L^2 - R^2} + 4A^2 m_i^2},$$

where $L^2 \equiv (m_{\tilde{\nu}}^2)_{ll} + (T_3 - e \sin^2 \theta_W)m_{\tilde{\nu}}^2 \cos 2\beta$, $R^2 \equiv (m_{\tilde{\nu}}^2)_{ll} + (e \sin^2 \theta_W)m_{\tilde{\nu}}^2 \cos 2\beta$, with $T_3 = -1/2$ and $e = -1$ for the down-type charged sleptons, and the effective trilinear scalar coupling term is denoted as $A \equiv (A_E)_{ll} - \mu \tan \beta$. $M_{\mu}$ and $M_{\nu}$ are slepton mass eigenstates obtained by diagonalising the slepton mass squared matrix. The trilinear $R'_p$ operator under consideration, i.e., $\lambda_{322}$ gives

\[\text{FIG. 2: Trilinear } R'_p \text{ violating contribution to neutrino masses.}\]
mass to the (33) element of the neutrino mass matrix. As a result, Eq. (11) can be simplified to

\[
(m_{\nu})_{33} \simeq \frac{1}{16\pi^2} |\lambda_{322}|^2 \frac{A_m^2}{\sqrt{(L^2 - R^2)^2 + 4A^2m_1^2}} \ln \left( \frac{M_{p1}^2}{M_{p2}^2} \right).
\]  

(13)

Considering the central values for the neutrino mass squared and mixing parameters [53] (with the choice, CP violating phase $\delta = 0$), the central value of the 33 element of the neutrino mass matrix for normal and inverted hierarchy becomes

\[
(m_{\nu})^{NH}_{33} = 0.023 \text{ eV}, \\
(m_{\nu})^{IH}_{33} = 0.031 \text{ eV}
\]  

(14)

From Eq. (13), it is straightforward to show that in the limit $(R^2 \gg L^2) \equiv \bar{m}^2 \gg A^2$, the same equation gives the following bound on the $A$ parameter as

\[
A \ln \left( \frac{M_{p1}}{M_{p2}} \right) \leq \mathcal{O}(10) \text{ GeV},
\]  

(15)

for $\mathcal{O}(1) \lambda_{322}$ and $\bar{m}$, i.e., the first two generations of right sleptons masses are in the ballpark of $\mathcal{O}(10 \text{ TeV})$. Therefore, we observe that in order to consider $\lambda_{322} \sim \mathcal{O}(1)$, one needs to satisfy Eq. (15)\(^4\) which invokes a cancellation between the soft SUSY breaking $A$ term in the charged slepton sector and the $\mu$ term.

IV. NUMERICAL ANALYSIS AND BENCHMARKS

From the previous discussions, it is clear that only $\lambda_{322}$ plays dominating role in ameliorating the tension between the observed muon anomalous magnetic moment and the SM expectation. In the limit when only $\lambda_{322}$ is non-zero, whereas all the other trilinear $R_p$ violating couplings are vanishingly small, Eq. (10) further simplifies to [16]

\[
\left[ a_{\mu}^{\lambda} \right]_{\text{RPV}} \simeq 34.9 \times 10^{-10} \left( \frac{100 \text{ GeV}}{\bar{m}} \right)^2 |\lambda_{322}|^2. 
\]  

(16)

It is conspicuous that an order one coupling can explain the muon anomalous magnetic moment event within 1σ of the central measured value.

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\(^{4}\) The issues pertaining to neutrino masses and muon $(g - 2)$ anomaly in the framework of RPV SUSY can also be found in Refs. [54, 55].

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\(^{5}\) In order satisfy the muon $(g - 2)$ and the LEP bound on the sneutrino mass simultaneously leads to $\lambda_{322} \geq 0.5$, however we restrict ourselves within $\lambda_{322} \leq 1.2$. 

In order to have a complete and concrete picture, we use the trilinear R-parity violating model implemented in SARAH v-4.4.6 [56, 57]. The spectrum has been generated using SPheno v-3.3.3 [58, 59]. FlavorKit [60] is used to ensure that the benchmark points are consistent with all relevant flavour violating observations. We fix the following parameters, such as, the bino mass parameter $M_1 = 300$ GeV, the wino mass parameter $M_2 = 1.7$ TeV, the Higgsino mass term $\mu = 200$ GeV, the gluino mass $M_3 = 1.5$ TeV, $\tan \beta = v_u/v_d = 20$ and the mass of the CP-odd Higgs $M_A = 400$ GeV. $\lambda_{322}$ is varied from 0.5 to 1.2 keeping all other $R_p$ couplings to zero\(^5\). We also vary the soft mass squared term of the slepton doublet in the limit $3 \times 10^4 \text{ GeV}^2 \leq (m_{\tilde{\tau}_1}^2)^2 \leq 2.5 \times 10^5 \text{ GeV}^2$ and chalk out the parameter space by putting the $\Delta a_\mu$ constraints within 1 & 2σ regime.

In Fig. 3 we show the 1 and 2σ constraints on $\Delta a_\mu$ in the $m_{\tilde{\nu}_1} - \lambda_{322}$ plane where $m_{\tilde{\nu}_1} \equiv m_{\tilde{\nu}_\mu}$. 

FIG. 3: 1 and $2\sigma$ limits on $\Delta a_\mu$ are shown in red and yellow colours respectively in the $m_{\tilde{\nu}_1} - \lambda_{322}$ plane where $m_{\tilde{\nu}_1} \equiv m_{\tilde{\nu}_\mu}$. 

1
our case. The present lower bound on $\tilde{\mu}$ stands at 300 GeV [12, 13], and thus this bound can also be mapped to an lower bound on $\tilde{\nu}_\tau$ mass. However, it is also important to note that by reducing the branching ratio of $\tilde{\nu}_\tau \to \mu^+\mu^-$, one can relax the bound considerably. For example, we check that for BR($\tilde{\nu}_\tau \to \mu^+\mu^-$) \sim 70\%, the bound on the sneutrino mass reduces to 250 GeV, while for BR($\tilde{\nu}_\tau \to \mu^+\mu^-$) \sim 50\% the bound on the same is around 220 GeV.

In our scenario, the partial decay width of the sneutrino (in this case the tau sneutrino) decaying to $\ell^+\ell^-$ is given by

$$\Gamma(\tilde{\nu}_i \to \ell^+_j\ell^-_k) \simeq \frac{1}{16\pi} \lambda_{ijk}^2 m_{\tilde{\nu}_i}.$$  (17)

Further, if kinematically allowed, the sneutrino can also undergo a two body decay with a tau neutrino and a neutralino or a tau lepton associated with a chargino in the final state. The neutralino and chargino would also undergo a three body decay in the RPV framework. The two body decay widths of the sneutrino are noted below

$$\Gamma(\tilde{\nu} \to \tilde{\chi}^0_2\nu) = \frac{g^2 |V_{11}|^2 m_{\tilde{\nu}}}{32\pi \cos^2 \theta_W} B(m_{\tilde{\chi}^0_2}^2/m_{\tilde{\nu}}^2),$$

$$\Gamma(\tilde{\nu} \to \tilde{\chi}^+\ell^-) = \frac{g^2 |V_{11}|^2 m_{\tilde{\nu}}}{16\pi} B(m_{\tilde{\chi}^+}^2/m_{\tilde{\nu}}^2).$$  (18)

where $V_{11}$ is one of the mixing matrix elements in the chargino sector and $Z_{ij}$ is the neutralino mixing matrix element. The $B$ function is defined as $B(x) = (1 - x)^2$. In the presence of large $\lambda_{322}$, which is also motivated from the perspective of fitting $\Delta a_\mu$, the partial decay width of the sneutrino decaying to a pair of leptons will dominate over the other decay modes. In Fig. 4, we portray the branching ratios of the lightest sneutrino as a function of its mass. During this scan, we ensure that all the points satisfy the Higgs mass and branching ratio constraints and also the low energy experimental constraints. In addition, care has been taken in removing all the tachyonic states from the scan. The points are also consistent within 2\sigma error of the $\Delta a_\mu$ parameter. All the parameters are fixed at the previously mentioned values expect for $M_1$. In the first column of Fig. 4, $M_1$ is fixed at 300 GeV, while in the lower panel $M_1 = 10$ GeV. As a result, sneutrino decays to charginos with associated leptons and neutrino+neutralino final states are highly suppressed due to phase space consideration. However, the bino-like neutralino mass parameter can be light (we choose it to be 10 GeV). There are two major constraints for light bino-like neutralino, for example, one has to check if Higgs partial decay width into this channel is satisfied or not and secondly, in the RPV scenario, the light neutralino can decay into final states involving fermions and can avoid constraints from its overproduction in the early universe [62, 63]. In addition, the added advantage of this scenario is the presence of light neutralino opens up new decay modes of the sneutrino. Thus the effective branching ratio of this sneutrino decaying to two muon final state can be reduced. As a result, the branching ratio of the stau decaying to $\mu\mu_\tau$ also reduces and thus relaxes the bound on the left handed stau mass. This in turn also implies the left-chiral sneutrino mass bounds can be relaxed further as both the masses are controlled by the same parameter. Before we proceed any further, we let us give a brief outline of the search for heavy di-muon resonances at the Tevatron and LHC experiments.

Heavy resonances decaying to a pair of muons naturally comes in many extensions of the SM with additional gauge groups. Both the ATLAS and CMS collaborations have searched for the heavy spin-1 resonance $Z'$ via di-muon final states at the 7 and 8 TeV run of LHC [64–66]. Non-observation of signatures of the signal events leads to the 95% C.L. upper limits on the production cross-section times branching ratios over a range of di-muon invariant masses. In Fig. 5, the black dotted and red dashed lines indicate the corresponding limits obtained from the ATLAS and CMS collaborations by the LHC-8 data respectively. Moreover, the CDF collaboration at the Tevatron experiment has also performed a study of di-muon resonances from the direct production of a
sneutrino or $Z'$ with 1.96 TeV data [67, 68]. However, we do not consider the bounds coming from the sneutrino production since it involves the $\lambda'$ LQD coupling in the production process which we set to be zero throughout our analysis. We find that our di-muon resonances, shown in blue dashed double-dotted (13 TeV) and red solid (8 TeV) lines, have smaller cross-sections compared to the ATLAS and CMS limits as elaborated later in the text.

Furthermore, we also consider the present bounds obtained from exotic searches at the LHC with final state topologies similar to ours i.e., with four muons among which two are positively charged and two negatively charged [69]. The ATLAS collaboration has searched for doubly charged Higgs bosons decaying to a pair of same sign muons, thus giving rise to the same final state signature. We again translate the 95% C.L. upper limit on the production process which we set to be zero through our analysis. We find that our di-muon resonances, shown in blue dashed double-dotted (13 TeV) and red solid (8 TeV) lines, have smaller cross-sections compared to the ATLAS and CMS limits as elaborated later in the text.

It is to be noted that, since we set $\lambda'$ to zero, the sneutrinos are produced at the LHC via only Higgs boson and off-shell $Z$ mediation and cross-section naturally becomes much smaller compared to present upper bounds except the doubly charged Higgs boson search process. The bound on the di-muon mass and hence on the sneutrino mass stands roughly at 290 GeV, similar to what we obtain translating the LHC bounds on the sleptons from the direct searches. Keeping all these bounds in mind, in Table II, we show the benchmark points pertaining to two relevant scenarios under consideration. The parameters which are fixed are $\tan \beta = 20$, $\mu = 200$ GeV, $M_A = 400$ GeV, $M_2 = 1.7$ TeV, $M_3 = 1.5$ TeV, $A_t = -1.9$ TeV, $\lambda_{322} = 1.2$, $(m_{\tilde{\chi}_1}^2)_{33} = 8.92 \times 10^4$ (GeV)$^2$ and $1.1 \times 10^5$ (GeV)$^2$ respectively. $M_1$ is fixed at 10 GeV (BP1) and 300 GeV (BP2) respectively. Obtained $\Delta a_\mu$ is within 2$\sigma$ error bar of the central value.

![FIG. 5: Present 95% C.L. upper limits on the $\sigma \times BR$ for different values of heavy resonance masses.](image)

| Panel     | BP1 | BP2 |
|-----------|-----|-----|
| $m_\nu$ (GeV) | 124.2 | 124.3 |
| $M_{H^0}$ (GeV) | 413 | 410 |
| $M_{H^+}$ (GeV) | 421 | 418 |
| $m_{\tilde{\nu}}$ (GeV) | 1622 | 1622 |
| $m_{\tilde{\tau}}$ (GeV) | 835 | 835 |
| $m_{\tilde{\chi}_1}$ (GeV) | 291 | 323 |
| $\lambda_{111}$ (GeV) | 204 | 204 |
| $\lambda_{212}$ (GeV) | 1711 | 1711 |
| $\lambda_{313}$ (GeV) | 9 | 310 |
| $\lambda_{414}$ (GeV) | 206 | 208 |
| $\lambda_{515}$ (GeV) | 210 | 309 |
| $\lambda_{616}$ (GeV) | 1711 | 1711 |
| $BR(b \rightarrow s \gamma) \times 10^4$ | 2.57 | 2.57 |
| $BR(B_s \rightarrow \mu^+ \mu^-) \times 10^{10}$ | 3.96 | 3.98 |
| $\Delta a_\mu \times 10^{10}$ | 19.6 | 19.7 |

TABLE II: Mass spectrum and a few observables for the two benchmark points.

V. COLLIDER ANALYSIS

Search for the new physics signatures with multiple leptons has always been considered as the golden channel mostly due to the cleanliness of the final state topology. Both the ATLAS and CMS collaborations at the LHC have searched for new resonances via lepton-rich signatures in the context of R-parity violating MSSM [70–72]. However, the final state topologies that have been studied by the CDF collaborations at the Tevatron and ATLAS collaborations at the LHC includes heavy neutral particles decaying to $e\mu$, $e\tau$, or $\mu\tau$ [70, 73]. From their analysis, we find that the $e\mu$ channel provides the best sensitivity due to better resolution for the electrons and muons (see Fig.2 of Ref. [70]). In this paper, we study possibly the cleanest final state topology which contains four isolated muons, which comes from the pair production of lightest sneutrino ($\tilde{\nu}_1 \equiv \tilde{\nu}_\tau$) which subsequently decays to a pair of muons through a non-zero $\lambda_{322}$ R-parity violating coupling. We reiterate that this channel is also interesting from the perspective of $(g-2)_\mu$ anomaly.

We perform the collider analysis for the two benchmark points already introduced. We generate signal events using MadGraph (v5 2.2.2) [74] where the main sneutrino pair production channel involves $Z$ mediation. We then pass the events to PYTHIA (v 6.4.28) [75] for hadronization and showering with CTEQ6L1 [76] parton
density function. The final state of interest contains four isolated muons with no real source of missing energy. The possible SM backgrounds that can mimic the signal topology are as follows: (i) SM Higgs boson production via gluon-gluon fusion, vector boson fusion, associated production processes with $H \rightarrow ZZ^* \rightarrow 4\mu$ final state. (ii) Direct production of pair of SM gauge bosons i.e., $WW$, $WZ$ and $ZZ$ with $W/Z$ decay leptonically. (iii) $Z$-jets and $t\bar{t}$ processes\(^6\). Similar to the signal events, the background processes are also simulated using Madgraph and then passed to PYTHIA. After generating the signal and background events, we apply the following kinematic cuts, which are more-or-less in line with those applied in a similar analysis by ATLAS collaboration\(^7\). We select the events with four isolated muons with $p_T > 10$ GeV and $|\eta| < 2.5$. The isolation criteria imposed on the muons are (a) that the angular separation $\Delta R_{\ell\ell}$ between the lepton and jets\(^7\) should not be less than 0.4, and (b) that the sum of the scalar $p_T$ of all stable visible particles within a cone of radius $\Delta R = 0.2$ around the lepton should not exceed 10 GeV. In Fig. 6, we display the $p_T$ distribution of the leading isolated muon. Note that, for the signal events the leptons are relatively harder compared to SM backgrounds and this important feature can be used as a trigger of such events. For our signal events, the muons are coming from the ‘on-shell’ decay of the sneutrino ($\tilde{\nu}_1$) and thus one can reconstruct $\tilde{\nu}_1$ mass using the di-leptonic invariant mass. However, for processes like $ZH$, $WH$ with $H \rightarrow ZZ^*$, one $Z$ is on-shell while the other is off-shell, and thus di-muon invariant mass will have a long tail with a sharp peak at $M_Z \sim 91$ GeV. Among all possible di-muon invariant mass recombinations, the one with minimum mass difference $\Delta m = |m_{12} - m_{34}|$ is selected where $m_{12}$ and $m_{34}$ are two such di-muon invariant masses. We impose a $Z$ veto by requiring either of the di-muon invariant masses is greater than 100 GeV. Note that, by making such a choice we also reduce the contributions coming from processes like associated production of a $Z$ boson with $J/\psi$ and/or $\Upsilon$ significantly. In lower panel of Fig. 6, the di-muon invariant masses are shown for both the signal and background events, where for the signal events we show for two representative benchmark points BP1 and BP2 with masses $\sim 290$ GeV and $320$ GeV respectively. From the figure it is evident that a cut on the di-muon invariant mass $m_{\mu\mu} > 100$ GeV would help us to reduce the dominant SM backgrounds.

In Table III, we show the production cross-section ($\sigma_0$), the effective cross-section ($\sigma_{\mathrm{eff}} = \sigma_0 \times \epsilon$, with $\epsilon$ being the cut efficiency) for the two benchmark points BP1 and BP2 along with the three dominant SM backgrounds $WH$, $ZH$ and $ZZ$ with $H \rightarrow ZZ^*$. The cross-sections for the signal events are calculated using Madgraph at the leading order\(^8\), while we follow the LHC Higgs Cross Section Working Group report\[^79\] for the $WH$, $ZH$ backgrounds where they are calculated at NNLO QCD and

\[^7\] We reconstruct jets using FASTJET v3.1.0\[^77\] with anti-$k_T$ jet algorithm and jet radius $R=0.4$.

\[^8\] We use Prospino\[^78\] to calculate the $K$-factor associated to the slepton pair production process and find $K=1.2$ for slepton masses from 200 - 500 GeV.
TABLE III: Event summary for the signal and background events. The quantities $\sigma_0$ and $\sigma_{\text{eff}}$ represent the production cross-section and the effective cross-section respectively. The total cross-section is denoted by $\sigma_{\text{tot}}$. For the $WH$ and $ZH$ processes the Higgs boson is assumed to decay to $ZZ^*$ with $Z$ decaying to two muons. We calculate the signal significance $S = S/\sqrt{(S + B)}$ with $L = 100$ fb$^{-1}$ where $S$ and $B$ are total number of signal and background events.

| Process | $\sigma_0$ (fb) | $\sigma_{\text{eff}}$ (fb) | $\sigma_{\text{tot}}$ (fb) | Significance |
|--------|-----------------|-----------------|-----------------|--------------|
| $\tilde{\nu}_1\tilde{\nu}_1^*$ | 4.08 (BP1) | 1.512 | 1.512 | |
| | 2.64 (BP2) | 0.98 | 0.98 | |
| $WH$ | 1380 [79] | 0.014 | 1.716 | 8.4 (BP1) |
| $ZH$ | 868 [79] | 0.0022 | 5.9 (BP2) |
| $ZZ$ | 15000 [80] | 1.7 |

FIG. 7: Contour plot in the $\mathcal{L} - m_{\tilde{\nu}_1}$ plane for the bino-line neutralino mass parameter of $M_1 = 300$ GeV. Similar distribution can be obtained for $M_1 = 10$ GeV.

VI. CONCLUSIONS

We revisit the possibility of satisfying anomalous magnetic moment of the muon in the paradigm of R-parity violating MSSM. The relevant coupling, $\lambda_{322}$, which plays a major role in this process is identified. The low energy and neutrino mass constraints have been checked and can be rather easily satisfied even at the presence of an $O(1)$ value of this particular R-parity violating coupling. We show that this explanation of having large muon $(g-2)$ via R-parity violation can be tested directly at the LHC. An artifact of $O(1)$ $\lambda_{322}$ is the decay of the pair produced tau sneutrino in to a final state comprising of four muons. This is a so-called “golden channel” because of large signal efficiency and minuscule contribution from the SM backgrounds. We analyze all the relevant SM backgrounds and find that sneutrino masses upto 450 GeV can be probed with an integrated luminosity of 1000 fb$^{-1}$ at the 13 TeV LHC. Such a channel is yet to be investigated by both the ATLAS and CMS collaborations, and it is our hope that this work will motivate them to perform a dedicated analysis in this direction in near future.

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APPENDIX

In this appendix we elaborate the SUSY-RPC contribution to the anomalous magnetic moment of the muon [14, 27, 28]. In general the chargino-sneutrino loop dominates over the neutralino-smuon loop. We reiterate that when all the mass scales are of the same order, the chargino-sneutrino loop contribution shown in Eq.(2) reduces to the form...
The neutralino-smuon contribution can be written down under the same approximation as
\[ \delta \alpha_{\mu}^{\pm} = \frac{m_{\mu}}{16\pi^2} \left\{ \frac{m_{\mu}}{12M_S^2} \left( \frac{m_{\mu}^2}{2M_S^2 \cos^2 \beta} + g_2^2 \right) + \frac{2g_2^2}{3M_S \sqrt{2} M_S \cos \beta} \right\} \]
\[ = \frac{m_{\mu}^2}{192\pi^2 M_S^2} \left\{ \frac{g_2^2 + 4\sqrt{2} g_2^2}{\cos \beta} \right\} \]
\[ \simeq \frac{m_{\mu}^2 g_2^2}{192\pi^2 M_S^2} \left\{ 1 + \frac{6}{\cos \beta} \right\}, \quad (A.1) \]

In the large \( \tan \beta \) limit, Eq. (A.1) can be further simplified to
\[ \delta \alpha_{\mu}^{\pm} \simeq \frac{m_{\mu}^2 g_2^2}{192\pi^2 M_S^2} \frac{6}{\cos \beta} \]
\[ \simeq \frac{m_{\mu}^2 g_2^2}{32\pi^2 M_S^2} \tan \beta. \quad (A.2) \]

Similarly, the neutralino-smuon contribution can be written down under the same approximation as
\[ \delta \alpha_{\mu}^{0} = \frac{m_{\mu}^2}{192\pi^2 M_S^2} \left( g_2^2 - g_1^2 \right) \tan \beta. \quad (A.3) \]

Therefore, the total RPC-SUSY contribution converges to the form given in Eq. (5). An interesting point to note is that although the one-loop contributions \( a_{\mu}^{\pi, \pm} \) has a term linear in \( m_{\chi^0, \pm} \) (see Eq. (1) and (2)) but they are not enhanced by \( m_{\chi^0, \pm} \) as compared to the other terms [14]. The reason being, these terms also involve either a factor of \( y_{\mu} \) or \( X_{m_1}X_{m_2} \), which is again proportional to \( (M_Z^2)^{1/2} \) and therefore to \( y_{\mu} \). Hence, all the RPC contributions to \( (g - 2)_{\mu} \) are of the order of \( m_{\mu}^2 / M_S^2 \) as shown explicitly. On the other hand, for \( O(1) \) RPV \( \lambda \) type couplings the contribution to \( a_{\mu} \) also turns out to be of the same order as its RPC counterpart.

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