Limits on R–parity non–conservation in MSSM with gauge mediated supersymmetry breaking from the neutrinoless double beta decay

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The R-parity non–conservation phenomenology within the minimal supersymmetric standard model with gauge mediated supersymmetry breaking has been studied. New constraints on the lepton number violating constant \( \lambda_{111} \) were imposed by non-observability of the neutrinoless double beta decay. We find that \( \lambda_{111} \) depends strongly on the effective supersymmetry breaking scale \( \Lambda \) only and deduce limits which are much stronger than those previously found in literature.

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Gauge–mediated theories of supersymmetry (SUSY) breaking (GMSB) have attracted a great deal of attention recently because of their high predictability, a natural solution of the flavour problem, a new phenomenology and much less free parameters compared to the Minimal Supersymmetric Standard Model (MSSM), where SUSY breaking is mediated by the gravitational interaction \( 111 \). In the GMSB models SUSY breaking is transmitted to the superpartners of quarks, leptons and gauge bosons via the usual \( SU(3) \times SU(2) \times U(1) \) gauge interactions and occurs at the scale of the order \( M_{SUSY} \sim 10^{10} \text{ GeV} \). Gauginos and sfermions acquire their masses through interactions with the messenger sector at the one– and two–loop levels respectively, resulting in different phenomenology from the MSSM one of the low–energy world. In these models flavour–diagonal sfermions mass matrices are induced in a rather low energy scale, therefore they supply us with a very natural mechanism of suppressing flavour changing neutral currents (FCNC). Moreover, since the soft masses arise as gauge charges squared, the sizeable hierarchy proportional to the gauge quantum numbers appears among the superpartner masses. As in MSSM the electroweak symmetry breaking (EWSB) is also driven by negative radiative corrections to the up–type Higgs mass resulting from the large top–quark Yukawa coupling and stop masses. Thus, with EWSB constraints the minimal GMSB models are highly constrained and strongly predictive. Recently renewed interest in GMSB \( 111 \) is then understood.

Motivated by the above features of the GMSB models we study the R–parity breaking phenomenology of MSSM and propose to use non–standard processes like the neutrinoless double beta decay \( 0\nu\beta\beta \) for deducing the limits on some R-parity breaking constants. In the previous studies such estimates were performed in the framework of MSSM with supergravity mediated SUSY breaking by means of additional assumptions relating sfermion and gauginos masses \( 111 \) or using the GUT constraints \( 111 \). Within the GMSB models one can find quantitatively new constraints, which may serve as a hint for further accelerator searches of R–parity breaking.

The R–parity imposed on MSSM can be explicitly violated by the trilinear \( 111 \) and bilinear \( 111 \) terms in the superpotential. The trilinear terms lead to the lepton number and flavour violation, while the bilinear terms generate the non–zero vacuum expectation values for the sneutrino fields \( \tilde{\nu}_i \), causing neutrino–neutralino mixing and electron–chargino mixing. Due to the lepton number violation, supersymmetric models with the R–parity broken may serve as a basis for description of some exotic processes. In the case of non–observability of \( 0\nu\beta\beta \) one can impose limits on the R–parity breaking parameters. It has been found that \( 0\nu\beta\beta \) is a powerful tool for such constraints \( 111 \).

To obtain the effective Lagrangian of the \( 0\nu\beta\beta \) decay we have to write the complete superpotential including both R–parity conserving and breaking parts:

\[
W = W_0 + W_R
\]

with

\[
W_0 = h^U_{ij} \tilde{Q}_i \tilde{H}_u \tilde{u}^c_j + h^D_{ij} \tilde{Q}_i \tilde{H}_d \tilde{d}^c_j + h^L_{ij} \tilde{L}_i \tilde{H}_u \tilde{\nu}^c_j + \mu \tilde{H}_d \tilde{H}_u.
\]

and

\[
W_R = \lambda_{ijk} \tilde{L}_i \tilde{L}_j \tilde{\nu}^c_k + \lambda'_{ijk} \tilde{Q}_i \tilde{L}_j \tilde{\nu}^c_k + \lambda''_{ijk} \tilde{\nu}^c_i \tilde{d}^c_j \tilde{d}^c_k.
\]

In Eqs. \( 111 \) \( \tilde{Q}, \tilde{L} \) denote the quark and lepton \( SU(2) \) doublet superfields and \( \tilde{\nu}^c, \tilde{d}^c, \tilde{e}^c \) are corresponding \( SU(2) \) singlets. The Higgs superfields \( \tilde{H}_u, \tilde{H}_d \) contain scalar components giving mass to the up– and down–type quarks and leptons. In the R–parity breaking part we set \( \lambda_{ijk} = \lambda'_{ijk} = 0 \) to avoid the unsuppressed proton decay. Since supersymmetry in the low–energy world is broken, the Lagrangian of the theory is supplemented with the "soft" supersymmetry breaking terms:
\[ -\mathcal{L}_{soft} = \left( A^U_{(ij)} h^U_i \bar{Q}_i H_u \bar{u}^c_j + A^D_{(ij)} h^D_i \bar{Q}_i H_d \bar{d}^c_j + A^E_{(ij)} h^E_i \bar{L}_i H_d \bar{e}^c_j + h.c. \right) \\
+ B_{\mu}(H_d H_u + h.c.) + m^2_{H_u}|H_d|^2 + m^2_{H_u}|H_u|^2 \\
+ m^2_L |\bar{L}^c|^2 + m^2_e |\bar{e}^c|^2 + m^2_{\tilde{Q}} |\tilde{Q}|^2 + m^2_{\tilde{W}} |\tilde{W}|^2 + m^2_{\tilde{E}} |\tilde{E}|^2 \\
+ \left( \frac{1}{2} M_1 \bar{\psi}_B \psi_B + \frac{1}{2} M_2 \bar{\psi}_W \psi_W + \frac{1}{2} m_\eta \bar{\psi}_\eta \eta + h.c. \right). \]

where a tilde denotes the scalar partners of quark and lepton fields, while \( \psi_i \)'s are the spin-\( \frac{1}{2} \) partners of the gauge bosons.

Using the R-breaking term \( W_R \) one can find the lepton number violating Lagrangian.

\[ \mathcal{L}_{\chi_{11}^2} = -\mathcal{L}_{\chi_{11}^2} (\bar{u}_L, \bar{d}_L) \left( \begin{array}{c} e^c_R \\ \eta_R \\ e^c_L \\ \eta_L \end{array} \right) \]

Together with Lagrangians describing interactions among gluinos, neutralinos, fermions and sfermions we end up, after integration out of heavy degrees of freedom, with the final effective Lagrangian

\[ \mathcal{L}_{\chi_{11}^2} = \frac{G_{\mu}}{2m_p} \bar{\chi}_R \left[ \eta_{PS} J_{PS} J_{PS} - \frac{1}{4} \eta_T J_{\mu\nu} J_{\mu\nu} \right], \]

where the color singlet hadronic currents are \( J_{PS} = \bar{u} c d + \bar{u} s d_s, J_{\mu\nu} = \bar{u} c \sigma_{\mu\nu} (1 + \gamma_5) d, \) with \( \alpha \) as a color index and \( \sigma_{\mu\nu} = (i/2) [\gamma^\mu, \gamma^\nu] \). The effective lepton–number violating parameters \( \eta_{PS} \) and \( \eta_T \) in Eq. (6) accumulate fundamental parameters of MSSM:

\[ \eta_{PS} = \eta_{\chi \bar{e}} + \eta_{\chi f} + \eta_{\chi} + \eta_{\chi} + 7\eta_{\bar{g}}, \]

\[ \eta_T = \eta_{\chi} - \eta_{\chi \bar{f}} + \eta_{\bar{g}} - \eta_{\bar{g}}, \]

where

\[ \eta_{\bar{g}} = \frac{\pi \alpha_s}{6} \frac{\lambda^{\mu}_{11}}{G_F^2 m_\tilde{g}} \frac{m_p}{m_\tilde{g}} \left[ 1 + \left( \frac{m_\tilde{d}}{m_\tilde{g}} \right)^4 \right], \]

\[ \eta_{\chi \bar{e}} = \frac{2\pi \alpha_s}{G_F^2 m_\tilde{g}} \frac{\lambda^{\mu}_{11}}{m_\tilde{e}} \sum_{i=1}^4 \frac{m_\tilde{d}}{m_{\chi}} \left( m_\tilde{d} \right)^4 \]

\[ \eta_{\chi f} = \frac{2\pi \alpha_s}{G_F^2 m_\tilde{g}} \frac{\lambda^{\mu}_{11}}{m_\tilde{e}/m_\tilde{g}} \sum_{i=1}^4 \frac{m_\tilde{d}}{m_{\chi}} \left( m_\tilde{d} \right)^2 \]

So, the effective Lagrangian depends on many supersymmetric parameters. In order to find their low–energy values in the GMSB model, we proceed as follows.

The minimal model of GMSB considered in this paper consists of messenger fields which transform as single flavor of \( 5 + 5 \) of \( SU(5) \), i.e. they are \( SU(2)_L \), doublets \( l \) and \( \bar{l} \) and \( SU(3) \) triplets \( q \) and \( \bar{q} \). The \( SU(3) \times SU(2) \times U(1) \) gauge interactions of those fields communicate supersymmetry breaking from a hidden sector to the fields of the visible world via coupling to a gauge singlet \( S \) in the superpotential

\[ W = \lambda_3 S \bar{L} \bar{L} + \lambda_3 S \bar{q} \bar{q}. \]
The lowest and $F$–components of the singlet superfield $S$ acquire v.e.v.’s and set the overall scale for the messenger sector and SUSY breaking respectively. For $F \neq 0$ the messenger sector loses its supersymmetry:

$$m_b = M \sqrt{1 \pm \frac{\Lambda}{M}},$$

$$m_f = M,$$

where $M = \lambda S$, $\Lambda = \frac{E}{F}$.

The messenger fields transmit the SUSY breaking to the visible sector through loops containing insertions of $S$ and results with the gaugino and scalar masses:

$$m_{\lambda_i}(M) = N_m g \left( \frac{\Lambda}{M} \right) \frac{\lambda_i(M)}{4\pi} \Lambda,$$

$$m^2(M) = 2 N_m f \left( \frac{\Lambda}{M} \right) \sum_{i=1}^{3} k_i \left( \frac{\lambda_i(M)}{4\pi} \right)^2 \lambda^2.$$

In Eq. (14) we sum up over $SU(3) \times SU(2)_L \times U(1)_Y$ with $k_1 = \frac{2}{3} (\frac{1}{2})^2$, $Y = 2(Q - T_3)$, $k_2 = \frac{1}{2}$ for $SU(2)_L$ doublets and zero for singlets, and $k_3 = \frac{4}{3}$ for $SU(3)_C$ triplets and zero for singlets. $N_m$ is the number of messenger generations equal to 1 in the minimal model. The messenger threshold functions $g(x)$ and $f(x)$ are:

$$g(x) = \frac{1+x}{x^2} \log(1+x) + (x \to -x),$$

$$f(x) = g(x) - \frac{2(1 + x)}{x^2} [Li_2 \left( \frac{x}{1+x} \right) - \frac{1}{4} Li_2 \left( \frac{x}{1+x} \right)] + (x \to -x).$$

It is worthwhile to note, that $g(x) \simeq f(x) \simeq 1$ unless $M$ is close to $\Lambda$ and both functions go to an infinity as $M$ equals to $\Lambda$.

The formulae for gaugino and sfermion masses at the messenger scale $M$ set up the boundary conditions for renormalization group equations (RGE) used to calculate the low–energy spectrum. In our procedure (see also [14]) we firstly evolve all the gauge and Yukawa couplings up to the scale $M$ using the 2–loop standard model RGE’s below

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...In our procedure (see also [14]) we firstly evolve all the gauge and Yukawa couplings up to the scale $M$ using the 2–loop standard model RGE’s below the mass threshold for SUSY particles (initially set up to $M_{SUSY} = 1$ TeV) and MSSM RGE’s above that scale. When the couplings reach the scale M we find the values of gaugino and scalar masses and perform the RGE evolution of all the quantities back to $m_Z$. It is well known, that running of $m^2_{H_u}$ is dominated by negative contribution from the top Yukawa coupling, which drives this parameter to a negative value at some scale and causes dynamic breaking of the electroweak symmetry (EWSB). This mechanism additionally allows to express some GUT–scale free parameters in terms of low–energy ones.

The tree–level Higgs potential has the form:

$$V_0 = m_{H_u}^2 |H_u^0|^2 + m_{H_d}^2 |H_d^0|^2 + m_3^2 (H_d^0 H_u^0 + h.c.) + \frac{g_1^2 + g_2^2}{8} (|H_d^0|^2 - |H_u^0|^2)^2,$$

where $m_{H_u}^2 \equiv m_{H_{u,\mu}}^2 + \mu^2$, $m_{H_d}^2 \equiv B\mu$, and phases of fields are chosen so that $m_3^2 < 0$. Due to the Q–scale dependence of soft Higgs parameters, $V_0$ depends strongly on Q. Thus, minimization of $V_0$ can lead to different v.e.v.’s ($v_u, v_d$) for different choices of Q. This is the reason why it is much more convenient to minimize the full one–loop Higgs effective potential. The procedure leads to the set of two equations:

$$|\mu|^2 + \frac{m_3^2}{2} = \frac{(m_{H_d}^2 + \Sigma_d) - (m_{H_u}^2 + \Sigma_u) \tan^2 \beta}{\tan^2 \beta - 1},$$

$$\sin 2\beta = \frac{-2B\mu}{(m_{H_u}^2 + \Sigma_u) + (m_{H_d}^2 + \Sigma_d) + 2|\mu|^2},$$

where $\Sigma_u, \Sigma_d$ are given in e.g. [15]. In order to minimize the stop contribution to the finite corrections we equal the minimization scale $Q_{min}$ to the geometric mean of stop masses.

EWSB causes mixing among many particles. In particular, mixing in the gaugino sector results in four physical neutralinos $\chi_i (i = 1, 2, 3, 4)$,
\[ \chi_i = \sum_{j=1}^{4} N_{ij} \psi_j. \]  

(20)

In this equation the matrix \( N \) diagonalizing the matrix \( M_\chi \), is real and orthogonal. Thus, for real \( N_{ij} \) neutralino masses are either positive or negative. If necessary, a negative mass can always be made positive by a redefinition of the relevant mixing coefficients \( N_{ij} \rightarrow iN_{ij} \).

Similar mixing appears in the slepton and squark sector. After calculation of the scalar particles masses and obtaining tree level values for \( \mu \) and \( B\mu \), one is able to set up thresholds in RGE’s for gauge and Yukawa couplings. Repeating the procedure by running everything up to the M scale and setting the M–scale conditions for soft parameters again, we proceed back to the \( m_Z \) scale. At this point we calculate new mass eigenstates, run everything up to \( Q_{min} \), minimize the one–loop corrected Higgs potential and perform RGE’s run of all the quantities to the \( m_Z \) scale. We iterate the procedure to obtain stable values of \( \mu \) and \( B\mu \).

Following the above scheme one can calculate the low–energy values of supersymmetric parameters appearing in the effective Lagrangian (6). Some of them can be restricted using experimental data for the non-observability of exotic nuclear processes. For example limits on \( \lambda^\prime_{111} \) can be deduced from the non–observability of \( 0\nu\beta\beta \). In this case the half–life for the neutrinoless double beta decay regarding all possibilities of hadronization of quarks can be written in the form:

\[
[T_{1/2}^{0\nu}(0^+ \rightarrow 0^+)]^{-1} = G_{01} \left| \eta_T M^N_q + (\eta_{PS} - \eta_T) M^N_f + \frac{3}{8}(\eta_T + \frac{5}{8}\eta_{PS}) M^{\pi N} \right|^2,
\]

where \( M^N_q, M^N_f \) and \( M^{\pi N} \) are defined in Ref. [10] and \( G_{01} \) is the standard phase space factor.

Using appropriate values for the nuclear matrix elements \( M^N_q, M^N_f, M^{\pi N} \) and the one–pion and two–pion contribution which dominate 2–N mechanism [10] we could obtain the neutrinoless double beta decay half–life as a function of a few supersymmetric parameters entering formula (6) only. Our procedure limits the number of these free parameters to: \( \Lambda, M, \tan \beta, sign(\mu) \) and \( N_m \). As the loop diagrams with messenger fields do not affect the A–terms considerably, we can equal the common soft SUSY breaking parameter \( A_0 \) to zero at the M–scale.

Setting the upper limit on the half–life of the neutrinoless double beta decay in \(^{76}\text{Ge}\) to \( 1.1 \times 10^{25} \text{y} \) [16] we are able to study constraints imposed by experiment on the R–parity breaking parameter \( \lambda^\prime_{111} \) as a function of the above mentioned free parameters.

We have found extremely weak dependence of \( \lambda^\prime_{111} \) on \( M, \tan \beta \) and \( sign(\mu) \) and much more pronounced influence of the \( \Lambda \) parameter. For the quantities commonly used in literature: \( \lambda^\prime_{111} \), \( (m_{\tilde{q}}/100\text{GeV}) \) and \( \frac{\sqrt{\lambda^\prime_{111}}}{(m_{\tilde{q}}/100\text{GeV})(m_{\tilde{g}}/100\text{GeV})^{1/2}} \) such a dependence is shown in Fig. 1 (a)-(c). The nuclear matrix elements used in these calculations were obtained for the bag model. It is worth noting that the strongest constraints come from the minimal scenario of GMSB \( (N_m = 1) \).
FIG. 1. Dependence of different representations of the $\lambda'_{111}$ parameter on $\Lambda$. Three curves for different numbers of the messenger multiplets ($N_m = 1, 2, 3$) are shown. Other free parameters are fixed as follows: $M = 100$ TeV, $\tan \beta = 20$, $\text{sign}(\mu) = +1$. For details see the text.

In conclusions, we have applied the experimental limits on the neutrinoless double beta decay of $^{76}$Ge to analysis of the R–parity non–conservation phenomenology within the GMSB theory. We have shown that this process imposes strong limits on the lepton number violating constant $\lambda'_{111}$. We have also found that $\lambda'_{111}$ depends significantly on $\Lambda$ only. Moreover, this dependence is stabilizing in the case of $\frac{\lambda'_{111}}{(m_q/100 \text{GeV})^2 (m_\nu/100 \text{GeV})^{1/2}}$ for $\Lambda > 70$ TeV which allows to estimate the lepton number violating constant:
\[
\frac{\lambda'_{111}}{(m_q/100\text{GeV})^2(m_g/100\text{GeV})^{1/2}} < 3.2 \times 10^{-5}.
\] (22)

At this point we would like to compare our results with the constraints on the R–parity broken MSSM coming from other processes. The limits obtained from considering the charged current universality violation impose the following bound: \(\lambda'_{111} \leq 0.03\frac{m_q}{100\text{GeV}}\) [17]. The above limit has been found within the gravity mediated MSSM. One can observe that in the GMSB case constraints are almost 2 orders of magnitude stronger.

Comparison with other estimates [9,11] shows that limit (22) is approximately one order of magnitude stronger. These results may be helpful for planning future searches for R-parity breaking signatures in different experiments.

The presented approach serves as a consistent and free of ad hoc assumptions about mutual dependence of model parameters one. To the best our knowledge we are first applying the gauge mediated supersymmetry breaking mechanism to constrain the R–parity breaking through the neutrinoless double beta decay.

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