Constraints on \( R \) parity and \( B \) violating couplings in gauge-mediated supersymmetry breaking models

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Abstract

We consider the proton decay involving a light gravitino or axino in gauge-mediated supersymmetry breaking models to derive constraints on the \( R \) parity and baryon number violating Yukawa couplings. Bounds on all nine coupling constants are obtained by considering the decay amplitudes at one-loop order.

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In supersymmetric models, there can be renormalizable gauge-invariant terms in the superpotential which violate the baryon number $B$ or the lepton number $L$. To avoid such terms, one usually introduces an additional discrete symmetry, the so called $R$ parity ($R_p = (-1)^{3B+L+2S}$). Although $R_p$ conservation leads to a consistent theory, there is no compelling theoretical reason to assume this symmetry. It is therefore an interesting possibility to have an explicit $R_p$ violation which may lead to interesting phenomenological consequences [1]. In the minimal supersymmetric standard model, the most general $R_p$-violating superpotential is given by

$$\frac{1}{2} \lambda_{ij k} L_i L_j E^c_k + \lambda'_{ij k} L_i Q_j D^c_k + \frac{1}{2} \lambda''_{ij k} U^c_i D^c_j D^c_k,$$

where $L_i$ and $Q_i$ are the $SU(2)$-doublet lepton and quark superfields and $E^c_i, U^c_i, D^c_i$ are the singlet superfields, respectively. Here $i,j,k$ are generation indices and we assume that possible bilinear terms $\mu_i L_i H_2$ are rotated away. Obviously the first and second terms in (1) violate $L$, while the third violates $B$. Since $\lambda_{ij k} = -\lambda_{jik}$ and $\lambda''_{ij k} = -\lambda''_{ikj}$, $R_p$ violations are described by 45 complex Yukawa couplings (9 in $\lambda$, 27 in $\lambda'$ and 9 in $\lambda''$).

It is well known that the consideration of proton decay provides a very stringent constraint on the product of $\lambda'$ and $\lambda''$:

$$|\lambda'_{112} \lambda''_{112} |, |\lambda'_{123} \lambda''_{113} | \leq 10^{-24},$$

where the squark masses are assumed to be around 1 TeV [2]. This bound has been obtained from a squark-mediated proton decay at tree level which does not involve heavy generation particles and thus applies for the particular combination of generation indices as is shown above. One may then expect that other products of $\lambda'$ and $\lambda''$ are allowed to be large. However it has been noted [3] that for any pair of $\lambda'$ and $\lambda''$, there is always at least one diagram relevant for the proton decay at one-loop level. This means that all products of $\lambda'$ and $\lambda''$ can be constrained by proton decay and a more detailed analysis leads to [3]

$$|\lambda' \cdot \lambda''| \leq 10^{-9},$$

for any pair of $\lambda'$ and $\lambda''$.

As was noted in [4,5], if there is a light fermion (lighter than the proton) which does not carry any lepton number, proton decay can be induced by a $B$ violating but $L$ conserving interaction alone, for instance by the $\lambda''$ couplings alone. Perhaps the most interesting class of models predicting such a light fermion are supersymmetric models in which supersymmetry (SUSY) breaking is mediated by gauge interactions [6]. In such models, the squark and/or gaugino masses, i.e. the soft masses in the supersymmetric standard model (SSM) sector, are given by $m_{soft} \simeq (\frac{\alpha_p}{\pi})^n \Lambda_S$ where $\Lambda_S$ corresponds to the scale of spontaneous SUSY breaking and the model-dependent integer $n$ counts the number of loops involved in transmitting SUSY breaking to the supersymmetric standard model sector. On the other hands, the gravitino mass is suppressed by the Planck scale $M_P \simeq 2 \times 10^{18}$ GeV, $m_{3/2} = \Lambda_S^2 / M_P$. Assuming that $m_{soft}$ is at the weak scale and taking $n = 1 \sim 3$ for instance, we have $m_{3/2} \simeq 10^{-1}$ eV $\sim 10$ MeV which is far below the proton mass.

Another interesting candidate for a light fermion without carrying any lepton number is the axino in supersymmetric models with a spontaneously broken global $U(1)_{PQ}$ symmetry.
If SUSY breaking is mediated by gauge interactions, the axino mass is given by \( m_\tilde{a} \approx (\alpha/\pi)^n \Lambda_5^2 / F_a \) where \( m \) is again a model-dependent (but typically not less than \( n \)) integer and \( F_a \) denotes the scale of spontaneous \( U(1)_{PQ} \) breaking \[1\]. Obviously in this case the axino can be lighter than the proton for a phenomenologically allowed \( F_a \geq 10^{10} \) GeV. In other type of models in which SUSY breaking is transmitted by supergravity interactions, the gravitino mass is fixed to be of the weak scale order, however there is still a room for an axino lighter than the proton \[9\]. As was pointed out in Ref. \[9\], some supergravity-mediated models lead to \( m_\tilde{a} \approx m_{3/2}/(m_{3/2}/M_P)^{1/2} \approx 1 \) keV for which the axino would be a good warm dark matter candidate \[10\].

In Ref. \[9\], proton decay involving a light gravitino or axino has been analysed at tree approximation to obtain a constraint on the \( R_p \) and \( B \) violating Yukawa coupling \( \lambda''_{ijk} \). Applying the naive dimensional analysis rule \[11\] for the hadronic matrix element of the effective 4-fermion operator induced by the tree diagram of Fig.1, the following stringent bounds on \( \lambda''_{112} \) (in the quark mass eigenstate basis) were obtained:

\[
\begin{align*}
\lambda''_{112} &\leq 5 \times 10^{-16} \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^2 \left( \frac{m_{3/2}}{1 \text{ eV}} \right), \\
\lambda''_{112} &\leq 7 \times 10^{-16} \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^2 \left( \frac{F_a}{10^{10} \text{ GeV}} \right) \left( \frac{1}{c_q} \right),
\end{align*}
\]

where \( \tilde{m} \) denotes the squark mass which is presumed to be universal. Here the dimensionless coefficient \( c_q \) describes the axino coupling to the light quarks \( q = (u, d, s) \) and is of order one for Dine-Fischler-Srednicki-Zhitnitski type axino, while it is of order \( 10^{-2} \sim 10^{-3} \) for hadronic-type axino. In this paper, we wish to extend the analysis of \[9\] by including one-loop effects and derive the constraints on the other components of \( \lambda'' \). The present upper bounds on \( \lambda'' \) are \( O(1) \) for sfermion mass \( \tilde{m} \sim 100 \) GeV except those on \( \lambda''_{112} \) and \( \lambda''_{113} \) \[12\]. As we will see, the derived upper bounds on \( \lambda''_{ijk} \) in gauge-mediated SUSY breaking models are much stronger than the currently existing bounds for a wide range of \( m_{3/2} \) and \( F_a \).

At low energy scales \( \sim 1 \) GeV where all massive particles are integrated out, the proton decay \( p \rightarrow \psi^+ \) light meson (\( \psi = \) gravitino or axino) can be described by an effective 4-fermion operator, \( \mathcal{O}_{\text{eff}} = \bar{uds}\psi \) or \( \bar{udd}\psi \), in the quark mass eigenstate basis. (See \[9\] for the detailed kinematic structure of these 4-fermion operators, which is not essential for our discussion in this paper.) At tree approximation, only \( \lambda''_{112} \) can produce such an effective operator (see Fig.1), thereby is constrained as \[9\]. In order for the other \( \lambda''_{ijk} \) to produce the flavor structure \( uds \) or \( udd \), it must be supplemented by flavor changing interactions in the model, which is possible at one-loop order. For instance, \( \lambda''_{212} \) and \( \lambda''_{312} \) can induce \( \mathcal{O}_{\text{eff}} \) once they are combined with the flavor change \( b \rightarrow d \), while \( \lambda''_{212} \) and \( \lambda''_{312} \) can do it with the flavor changes \( c \rightarrow u \) and \( t \rightarrow u \), respectively. The other four couplings need double flavor changes in order to lead to a proton decay, e.g. \( (c, b) \rightarrow (u, d) \) or \( (u, s) \) for \( \lambda''_{213} \) and \( \lambda''_{223} \), \( (t, b) \rightarrow (u, d) \) or \( (u, s) \) for \( \lambda''_{313} \) and \( \lambda''_{323} \).

To proceed, let us collect the couplings which are relevant for the proton decay into light gravitino or axino at one-loop order. First of all, one needs the following gravitino (\( G \)) \[13\] or axino (\( \tilde{a} \)) \[14\] couplings:

\[
\mathcal{L}_G = \frac{i}{4\sqrt{6}m_{3/2}M_P} \left[ \tilde{\lambda}^\alpha \gamma^\mu \sigma^\nu \partial_\mu GF_\mu^\nu + 2\sqrt{2}\psi_I (1 - \gamma_5) \gamma^\mu \gamma^\nu \partial_\mu GD_\nu \phi_I^* \right] + \text{H.c.}
\]
\[
\mathcal{L}_a = -\frac{c_I}{2F_a} \left[ i \partial_\mu \bar{\psi}_I \gamma^\mu (1 + \gamma_5) \tilde{a} \phi^*_I + \text{H.c.} \right] + \frac{c_\alpha}{32\sqrt{2}\pi^2 F_a} \left[ \tilde{\lambda}^\alpha \gamma^\mu (1 - \gamma_5) \tilde{a} F^\alpha_{\mu\nu} + \text{H.c.} \right],
\]

where \((\phi_I, \psi_I)\) and \((\lambda^\alpha, F^\alpha_{\mu\nu})\) stand for the chiral matter and gauge multiplets. Here the axino couplings \(c_\alpha\) to gauge multiplets are generically of order one, while the couplings \(c_I\) to matter multiplets are of order one only for matter multiplets carrying a nonzero \(-B^{(1)}\), for instance the following flavor changing interactions which we will use in the subsequent discussions include the charginos. The flavor changing interactions in \(R_p\)-conserving sector are highly suppressed. Then the necessary flavor change takes place through the exchange of \(W^\pm\), charged Higgs or charginos. The flavor changing interactions which we will use in the subsequent discussions include the \(W\)-boson coupling:

\[
- V_{ij} \frac{g}{\sqrt{2}} W_\mu^+ \bar{u}^i \gamma^\mu P_L d^j + \text{H.c.},
\]

and the charged Higgs boson coupling:

\[
V_{ij} H^+ \left[ f^{(d)}_j \tan \beta \bar{u}^i P_R d^j + f^{(u)}_i \cot \beta \bar{u}^i P_L d^j \right] + \text{H.c.},
\]

where \(g\) is the \(SU(2)\) gauge coupling, \(V_{ij}\) is the CKM matrix element, \(\tan \beta\) is the ratio of Higgs vacuum expectation values, and \(f^{(u,d)}_i\) denote the quark Yukawa couplings, i.e.

\[
f^{(u)}_i = \frac{g m^{(u)}_i}{\sqrt{2} m_W \sin \beta}, \quad f^{(d)}_i = \frac{g m^{(d)}_i}{\sqrt{2} m_W \cos \beta}.
\]

(Here \(m_W, m^{(u)}_i\) and \(m^{(d)}_i\) denote the masses of \(W\)-boson, up- and down-type quarks respectively.) In addition to these, we will use the following chargino-quark-squark interactions also:

\[
V_{ij} \left\{ \begin{array}{l}
\tilde{f}^{(u)}_i \sin \phi_L P_L + f^{(u)}_i \sin \phi_R P_R \bar{u}^i \tilde{d}^{i*}_L + f^{(d)}_j \sin \phi_L \tilde{\lambda}^+_1 P_L \bar{u}^i \tilde{d}^{i*}_R \\
+ \tilde{f}^{(u)}_i \sin \phi_L P_L + f^{(u)}_i \sin \phi_R P_R \bar{u}^i \tilde{d}^{i*}_L + f^{(d)}_j \sin \phi_L \tilde{\lambda}^+_2 P_L \bar{u}^i \tilde{d}^{i*}_R \\
+ \tilde{f}^{(u)}_i \sin \phi_L P_L + f^{(u)}_i \sin \phi_R P_R \bar{u}^i \tilde{d}^{i*}_L + f^{(d)}_j \sin \phi_R \tilde{\lambda}^+_2 P_L \bar{u}^i \tilde{d}^{i*}_R \\
+ \tilde{f}^{(u)}_i \sin \phi_L P_L + f^{(u)}_i \sin \phi_R P_R \bar{u}^i \tilde{d}^{i*}_L + f^{(d)}_j \sin \phi_R \tilde{\lambda}^+_2 P_L \bar{u}^i \tilde{d}^{i*}_R \\
+ \tilde{f}^{(u)}_i \sin \phi_L P_L + f^{(u)}_i \sin \phi_R P_R \bar{u}^i \tilde{d}^{i*}_L + f^{(d)}_j \sin \phi_R \tilde{\lambda}^+_2 P_L \bar{u}^i \tilde{d}^{i*}_R \end{array} \right\} + \text{H.c.},
\]

where \(\epsilon_R \equiv \text{sign}(\mu M_2 - m_W^2 \sin 2\beta)\) for the gaugino mass \(M_2\) and the Higgsino mass parameter \(\mu\). The chargino mixing angles \(\phi_{L,R}\) are given by

\[
\phi_{L,R} = \tan \theta_{12} \mp \frac{\tan \theta_{12}}{\sqrt{1 - \tan^2 \theta_{12}}}.
\]
\[\tan 2\phi_L = \frac{2\sqrt{2}m_W (M_2 \cos \beta + \mu \sin \beta)}{M_2^2 - \mu^2 - 2m_W^2 \cos 2\beta},\]
\[\tan 2\phi_R = \frac{2\sqrt{2}m_W (M_2 \sin \beta + \mu \cos \beta)}{M_2^2 - \mu^2 + 2m_W^2 \cos 2\beta}.\] (11)

All one loop diagrams which trigger a proton decay by having a \(\lambda''\)-vertex can be divided into the following three categories; (a) diagrams with radiative corrections to the \(\lambda''\)-vertex (Fig. 4), (b) box diagrams (Fig. 4), (c) diagrams with radiative corrections to the gravitino or axino vertex (Fig. 5). Relative to the tree diagram of Fig. 1, one loop diagrams involving \(\lambda''_{ijk}\) will be suppressed by the factor \(\xi_{ijk}\), more explicitly

\[\frac{A_{\text{loop}}^{(ijk)}}{\lambda''_{ijk}} = \frac{A_{\text{tree}}}{\lambda''_{112}},\] (12)

where \(A_{\text{tree}}\) denotes the tree amplitude of Fig. 1, while \(A_{\text{loop}}^{(ijk)}\) stand for the loop amplitudes of Fig. 4 and Fig. 5 which involve the insertion of \(\lambda''_{ijk}\). The upper bounds on \(\lambda''_{ijk}\) resulting from those one loop diagrams can be easily read off from (4) by taking into account the suppression factor \(\xi_{ijk}\):

\[\lambda''_{ijk} \leq 5 \times 10^{-16} \left( \frac{1}{\xi_{ijk}} \right) \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^2 \left( \frac{m_{3/2}}{1 \text{ eV}} \right),\]
\[\lambda''_{ijk} \leq 7 \times 10^{-16} \left( \frac{1}{\xi_{ijk}} \right) \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^2 \left( \frac{F_a}{10^{10} \text{ GeV}} \right) \left( \frac{1}{c_q} \right).\] (13)

In the following, we will estimate the size of \(\xi_{ijk}\) for the loop diagrams depicted in Fig. 4 and Fig. 5. Let us first consider the type (a) and (b) diagrams in Fig. 4. It turns out that type (a) diagrams (with the charged Higgs exchange) dominate in this case. The resulting suppression factors are given by

\[\xi_{ijp} \approx \frac{1}{(4\pi)^2 } f_{(u)}^{(1)} V_{iq} f_{(d)}^{(1)} V_{ij}^* = \frac{g^2}{16\pi^2} \frac{1}{m_W^2 \sin(2\beta)} m_{(u)}^{(1)} V_{iq} m_{(d)}^{(1)} V_{ij}^*,\] (14)

where \((p, q) = (1, 1), (1,2)\) or \((2,1)\). It is worth noting that these suppression factors are rather insensitive to the details of unknown superparticle masses.

Although it depends more on the details of superparticle spectrum, for \(\lambda''_{113}\) and \(\lambda''_{123}\), one can get a much stronger bound through the diagrams in Fig. 5. For instance, we find that the loop suppression factors of Fig. 5–(i) are given by

\[\xi_{123} \approx \frac{g^2}{16\pi^2} V_{31} V_{33}^{*} \frac{m_b}{m_W} \frac{m_t^2 \delta m_{t}^2}{\tilde{m}^4} \sin 2\phi_L \approx 5 \times 10^{-7} \left( \frac{\delta m_{t}^2}{\tilde{m}^2} \right) \left( \frac{300 \text{ GeV}}{\tilde{m}} \right)^2 \sin 2\phi_L,\]
\[\xi_{113} \approx \frac{g^2}{16\pi^2} V_{32} V_{33}^{*} \frac{m_b}{m_W} \frac{m_t^2 \delta m_{t}^2}{\tilde{m}^4} \sin 2\phi_L \approx 2 \times 10^{-6} \left( \frac{\delta m_{t}^2}{\tilde{m}^2} \right) \left( \frac{300 \text{ GeV}}{\tilde{m}} \right)^2 \sin 2\phi_L,\] (15)

where we have assumed that all superparticle masses including the chargino masses are the approximately same as the universal squark mass \(\tilde{m}\), and \(\delta m_{\tilde{\chi}}^2\) is the difference between the two chargino mass-squared.
$$\delta m^2_\chi \equiv |m^2_{\chi_1} - m^2_{\chi_2}|. \quad (16)$$

Here the extra \( m_t \)-dependence is due to the GIM-cancellation. In fact, one can consider a diagram which is similar to Fig.3–(i) but including the insertion of \( \lambda''_{212} \) or \( \lambda''_{312} \). However the amplitude of such diagram is heavily suppressed by the GIM mechanism, and thus it does not give a bound on \( \lambda''_{212} \) or \( \lambda''_{312} \) which would be stronger than the bound from Fig.3.

If the charginos are degenerate or the chargino mixing \( |\sin 2\phi_L| \ll 1 \), the bounds from Fig.3–(i) will be significantly weakened. In this case, the dominant contribution would come from Fig.3–(ii) or 3–(iii) which involves the insertion of the left-right squark mixing:

$$m_j^{(d)} \mu \tan \beta \tilde{d}_j^L \tilde{d}_R^i + m_i^{(u)} \mu \cot \beta \tilde{u}_L^i \tilde{u}_R^j + \text{H.c..} \quad (17)$$

The corresponding loop suppression factors are given by

$$\xi_{123} \approx \frac{g^2}{8\pi^2} V_{31} V_{33}^* \frac{\mu m_b}{m^2} \left( \frac{\mu}{\tilde{m}} \right) \left( \frac{300 \text{ GeV}}{\tilde{m}} \right)^3 \approx 4 \times 10^{-7},$$

$$\xi_{113} \approx \frac{g^2}{8\pi^2} V_{32} V_{33}^* \frac{\mu m_t}{m^2} \left( \frac{\mu}{\tilde{m}} \right) \left( \frac{300 \text{ GeV}}{\tilde{m}} \right)^3 \approx 1 \times 10^{-6}, \quad (18)$$

where again it is assumed that all superparticle masses are the approximately same as the universal squark mass \( \tilde{m} \).

In fact, for the axino case there arises an extra complication for Fig.3–(i) and Fig.3–(ii) since they involve the axino coupling to gauge multiplets \( c_{\alpha}/8\pi^2 \) in Eq. (3), while the tree diagram of Fig.1 involves only the axino coupling to the light quark multiplets \( c_I \) in Eq. (3) for \( I = q \) where \( q = (u, d, s) \) stands for the light quark multiplets. As a result, for the axino case, the correct suppression factors of Fig.3–(i) and Fig.3–(ii) are obtained by multiplying the factor \( c_{\alpha}/8\pi^2 c_q \) to the results of Eqs. (15) and (18). This point is irrelevant for the Dine-Fischler-Srednicki-Zhitnitskii type axino, however it may lead to one order of magnitude stronger bound for hadronic-type axino. In this paper we will ignore this extra complication for the sake of simplicity.

Applying the loop suppression factors of Eqs. (14), (15) and (18) to Eq. (13), one can easily derive the upper bounds on \( \lambda''_{ijk} \). We summarize the numerical results in Table I. In deriving these, we take \( \sin(2\beta) = 1 \) in (14) and ignored the contributions from Fig.3–(i) and Fig.3–(ii) for the axino case, which would lead to conservative results. We also assumed that all superparticle masses are the approximately same as the universal squark mass \( \tilde{m} \), and also \( \mu \approx \tilde{m} \). For the numerical values of the quark masses, CKM matrix elements and etc., the values in Ref. [15] are used.

To conclude, we have examined the proton decay involving a light gravitino or axino in gauge-mediated supersymmetry breaking models to derive constraints on the \( R \) parity and baryon number violating Yukawa couplings \( \lambda''_{ijk} \). Considering the decay amplitudes at one-loop order, we could get upper bounds on all of those couplings. The results summarized in Table I show that, for a wide range of the gravitino mass \( m_3/2 \), the bounds on all \( \lambda''_{ijk} \) are much stronger than the currently existing bounds. The bounds on \( \lambda''_{113} \) and \( \lambda''_{123} \) are particularly strong due to the contributions from Fig. 3. In supersymmetric models with \( U(1)_{PQ} \), if axino is lighter than the proton, all \( \lambda''_{ijk} \) are similarly constrained by the proton decay into light axino.
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TABLE I. Upper bounds on $\lambda''_{ijk}$ in gauge-mediated SUSY breaking models from the proton decay into light gravitino (bound I) or axino (bound II). Here the bounds on $\lambda''_{113}$ and $\lambda''_{123}$ are from Fig.4, while others from Fig.3. All superparticle masses are assumed to be the approximately same as the universal squark mass $\tilde{m}$, and $x_s \equiv \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)$, $x_{3/2} \equiv \left( \frac{m_{3/2}}{1 \text{ eV}} \right)$, $x_a \equiv \left( \frac{F_a}{10^{10} \text{ GeV}} \right) \left( \frac{1}{c_q} \right)$.

| Coupling   | Upper Bound I       | Upper Bound II      |
|------------|---------------------|---------------------|
| $\lambda''_{112}$ | $5 \times 10^{-16} x_s^2 x_{3/2}$ | $7 \times 10^{-16} x_s^2 x_a$ |
| $\lambda''_{113}$ | $3 \times 10^{-10} x_s^3 x_{3/2}$ | $7 \times 10^{-10} x_s^3 x_a$ |
| $\lambda''_{123}$ | $1 \times 10^{-9} x_s^3 x_{3/2}$ | $2 \times 10^{-9} x_s^3 x_a$ |
| $\lambda''_{212}$ | $3 \times 10^{-8} x_s^2 x_{3/2}$  | $4 \times 10^{-8} x_s^2 x_a$ |
| $\lambda''_{213}$ | $5 \times 10^{-8} x_s^2 x_{3/2}$  | $7 \times 10^{-8} x_s^2 x_a$ |
| $\lambda''_{223}$ | $3 \times 10^{-7} x_s^2 x_{3/2}$  | $4 \times 10^{-7} x_s^2 x_a$ |
| $\lambda''_{312}$ | $5 \times 10^{-9} x_s^2 x_{3/2}$  | $7 \times 10^{-9} x_s^2 x_a$ |
| $\lambda''_{313}$ | $1 \times 10^{-8} x_s^2 x_{3/2}$  | $1 \times 10^{-8} x_s^2 x_a$ |
| $\lambda''_{323}$ | $5 \times 10^{-8} x_s^2 x_{3/2}$  | $7 \times 10^{-8} x_s^2 x_a$ |
FIG. 1. Tree diagram for the proton decay into light gravitino or axino.
FIG. 2. One loop diagrams for the proton decay into light gravitino or axino.
FIG. 3. Other one loop diagrams relevant for $\lambda''_{113}$ and $\lambda''_{123}$. Here cross means the left-right squark mixing.