Large R-parity Violating Couplings and Grand Unification

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ABSTRACT

We consider the possibility that R-parity violating interactions of particles from second and third matter generations have large (up to 1) coupling constants, Λ. Such couplings have a number of phenomenological consequences: renormalization of $b - \tau$ mass ratio, generation of ν_τ mass in MeV region, etc. In Grand Unified models, where B- and L-violating couplings appear simultaneously, the proton decay can be forbidden in virtue of hierarchical flavour structure of Λ. However, due to Cabibbo-Kobayashi-Maskawa mixing, this decay is induced already in one-loop. Present experimental data give the upper bound $\Lambda \lesssim 3 \cdot 10^{-9}$ (or $|\lambda'\lambda''| \lesssim 5 \cdot 10^{-17}$, on the product of the L- and B-violating coupling constants, in more general context). The bounds can be avoided, if there is an asymmetry between the L- and B-violating couplings of usual matter fields. In the $SU(5)$ model the asymmetry can be related to the doublet-triplet splitting.

1 Introduction

The gauge invariance of the Standard Model and Supersymmetry [1] permits, besides usual Yukawa interactions,

$$
\begin{aligned}
W &= m_{E,i}/v_1 \; E_i^c (H_1^0 E_i - H_1^- \nu_i) \\
&+ m_{D,i}/v_1 \; D_i^{c\alpha} (H_1^0 D_i^\alpha - H_1^- U_j^\alpha V_{ji}^*) \\
&+ m_{U,i}/v_2 \; U_i^{c\alpha} (H_2^0 U_i^\alpha - H_2^+ V_{ij} D_j^\alpha) \\
&+ \mu(H_1^0 H_2^0 - H_1^- H_2^+),
\end{aligned}
\tag{1}
$$

also the couplings which violate either lepton or baryon number conservation [2]:

$$
\begin{aligned}
W_{\not{R}} &= \lambda_{ijk}(E_i \nu_j - \nu_i E_j)E_k^c \\
&+ \lambda'_{ijk} D_i^{c\alpha}(\nu_j V_{kl} D_l^\alpha - E_j U_k^\alpha) \\
&+ \lambda''_{ijk}\epsilon_{\alpha\beta\gamma} D_i^{c\alpha} D_j^{c\beta} U_k^{c\gamma}.
\end{aligned}
\tag{2}
$$

Here, $E_i^c, E_i, \nu_i, D_i^c, D_i, U_i^c, U_i$ are the superfields with charged leptons, neutrinos, down- and up-type-quarks; $i,j,k,l = 1,2,3$ are generation indices; $H_{1,2}^0$, H_1^-, H_2^+ are the Higgs supermultiplets, and $v_{1,2}$ are the vacuum expectation values of the scalar components of $H_{1,2}^0$. The supermultiplets W and $W_{\not{R}}$ are written in terms of superfields with fermion mass eigenstates, so that the Cabibbo-Kobayashi-Maskawa matrix V_{ij} appears in (1) and (2) explicitly; $m_{E,i}$, $m_{D,i}$, $m_{U,i}$ are the fermion masses. Another possible term in $W_{\not{R}}$, $\mu_i(\nu_i H_2^0 - E_i H_2^+)$, can be rotated away from the superpotential, by redefinition of the couplings in W and $W_{\not{R}}$.

A rich phenomenology can be related to the interactions (2). They result in B- or/and L-violating phenomena like $n - \bar{n}$ oscillations [3, 4, 5], proton decay [6, 7], generation of Majorana neutrino masses [8, 9]; they modify usual processes like μ-, β-decay [10], and lead to the decay of the lightest supersymmetric particle [11]. However, up to now no effects of (2) have been found which implies strong restrictions on the constants λ. In particular the proton decay searches allow to put the bound on certain couplings of lowest generations:

$$
|\lambda'\lambda''| \lesssim 10^{-24}
\tag{3}
$$

for squark masses around 1 TeV.

The smallness of at least some couplings (2) indicates that probably all the interactions (2) are absent in virtue of certain symmetry. Moreover, the absence of the terms (2)

ensures stability of the lightest supersymmetric particle which is considered as a favourite candidate for the cold dark matter. $W_{\not R}$ can be suppressed by R-parity or matter parity conservation. The corresponding symmetries may naturally follow from a class of grand unified symmetries like $SO(10)$ in models with minimal particle content. Alternatively B- or L-violating terms can be suppressed by symmetries which distinguish quarks and leptons.

In this paper we assume that R-parity (or some other symmetry which suppresses $W_{\not R}$) *is not exact* and the terms (2) are generated with sufficiently small coupling constants. In fact, the existing data strongly restrict the couplings of light generations, whereas the bounds on couplings of second and third generations are weak or absent (for latest discussion see [12]). In the same time it is natural to assume the hierarchy of constants λ [13]. Moreover, as the consequence of a horizontal symmetry, this hierarchy can be much stronger than that of the usual Yukawa couplings. Strong hierarchy of λ can be partially related to the fact that couplings in (2) involve three generation dependent fields, whereas Yukawa couplings contain only two such fields (see for latest discussion [12]). Thus the following pattern is possible: the constants λ for the first and second generations are very small and satisfy the existing bounds, while the couplings involving third generation particles are large and could be of the order 1.

Large R-parity violating couplings of third generation can manifest themselves in many ways.

At one-loop they induce the Majorana neutrino masses [8, 9]. They contribute to $K^0 - \bar{K}^0$ mixing, to the electric dipole of the neutron [4], to $Z \to b\bar{b}$ decay width [14], the decay of B meson $B^- \to K^0 K^-$ [12] etc..

Large λ's influence the running of usual Yukawa couplings. In particular, they modify the infrared fixed point of the top quark Yukawa coupling [15]. The restriction $\lambda''_{i33} < 0.4 - 0.5$ has been obtained from the condition that the top coupling does not blow up before the grand unification scale M_{GU}. Large B- or L-violating couplings of the heaviest generations can appreciably renormalize the $b - \tau$ mass ratio. It is shown [16] that for values $\lambda''_{233} = 0.15 - 0.30$ the $(b - \tau)$-mass unification at GU scale can be achieved for any value of $\tan \beta$ in the interval $2 - 50$.

The studies of the R-parity violation effects were performed mainly in the context

of Minimal Supersymmetric Standard Model. However remarkable convergency of the gauge couplings at the scale around $3 \cdot 10^{16}$ GeV [17, 18, 19] can be considered as strong indication of the supersymmetric unification of the strong and the electroweak interactions. Supersymmetry offers an elegant way to stabilize the gauge hierarchy, thus ensuring consistency of the picture. Moreover, the $b - \tau$ unification [20] can be achieved in the supersymmetric GU model only [21]. Note that λ-couplings, like the usual Yukawa couplings, will affect only weakly (at the two-loop level) the evolution of the gauge coupling constants. In this connection it is important to consider the properties and consequences of the interactions (2) in the GU theories. The first studies of R-parity violation in the context of Grand Unification have been performed in [22, 23, 24].

In this paper we consider the proton decay induced by R-parity violating couplings of heaviest (second and third) matter generations. We find new very strong bounds on λ in the $SU(5)$ with standard matter field content. The modifications of the model are discussed which allow us to get the asymmetry of B- and L-violating couplings and thus to avoid the bounds.

The paper is organized as follows. Properties of R-parity violating couplings in SUSY $SU(5)$ are discussed in sect. 2. We consider the proton decay induced by these couplings in sect. 3. The conditions are found at which the decay is forbidden in the lowest order of perturbation theory. However, being suppressed in lowest order, proton decay is inevitably generated by one-loop diagrams (sect. 4). The amplitudes of leading one-loop diagrams are estimated and the upper bounds on R-parity violating coupling constants are found. In sect. 5 we consider the generality of the bounds and the way to avoid them. Then (sect. 6) we discuss possible relations between asymmetry of the B- and L-violating interactions which allows one to avoid the bounds and the doublet-triplet splitting. Sect. 7. summarizes the results.

2 R-parity violating interactions in the $SU(5)$-supersymmetric model.

In the $SU(5)$ model one can introduce the following R-parity violating interactions [22]

$$\Lambda_{ijk}\bar{5}_i\bar{5}_j 10_k + \bar{5}_i(M_i + h_i\Phi)H, \tag{4}$$

where $i, j, k = 1, 2, 3$ are generation indices, Λ_{ijk} are the coupling constants and $\bar{5}_i$, 10_i are the matter superfields which can be written in terms of the standard model supermultiplets as:

$$\bar{5} = \begin{pmatrix} D^c \\ i\sigma_2 L \end{pmatrix} \qquad 10 = \begin{pmatrix} U^c & -Q \\ Q & -E^c i\sigma_2 \end{pmatrix}. \tag{5}$$

Here σ_2 is the Pauli matrix, $L = (\nu, E)$ and $Q = (U, D)$ are $SU(2)_L$ doublets. M_i are the mass parameters, h_i are the couplings, Φ and H are the 5-plet and 24-plet of Higgs fields.

Let us consider first the effects of Λ couplings, suggesting that the matter-Higgs mixing (second term in (4)) is negligibly small. The Λ_{ijk}-coupling (4) generate all the R-parity violating interactions (2). It is convenient to define Λ_{ijk} in the basis, where $SU(2)_L$-singlets u^c and d^c (fermionic components of U^c and D^c) coincide with mass eigenstates. This always can be done since u^c and d^c enter different $SU(5)$-multiplets. Note that due to the antisymmetry of 10-plets the interactions (4) are antisymmetric in generation indices: $\Lambda_{ijk} = -\Lambda_{jik}$.

Substituting multiplets (5) in (4) and comparing the resulting interactions with those in (2) we find the relations between original λ_{ijk} and Λ_{ijk} couplings at the GU scale:

$$\begin{aligned} \lambda_{ijk} &= \Lambda_{ijl} V_{lk} \\ \lambda'_{ijk} &= 2\Lambda_{ijk} \\ \lambda''_{ijk} &= \Lambda_{ijk}. \end{aligned} \tag{6}$$

As a consequence of quark and lepton unification in $SU(5)$, all types of R-parity violating couplings appear simultaneously. Moreover, different couplings λ, λ' and λ'' are determined by unique GU coupling Λ. As follows from (6), up to CKM matrix and factor 2 in λ' these couplings coincide at GU scale:

$$\lambda_{ijl} V_{lk}^{-1} = \frac{1}{2}\lambda'_{ijk} = \lambda''_{ijk}. \tag{7}$$

Evidently, there is no relative suppression of B- and L-violating couplings. Another feature of the Grand Unification is that L-violating couplings, λ'_{ijk}, should be antisymmetric in first two indices: $\lambda'_{ijk} = -\lambda'_{jik}$, similarly to other couplings. In the non-unified version (2) these couplings can have also a symmetric part.

The gauge coupling renormalization effects lead to modification of GU relations (6) at

the electroweak scale:

$$\lambda_{ijk} = 1.5 \ \Lambda_{ijl} V_{lk}$$

$$\lambda'_{ijk} = 2 \ (3.4 \pm 0.3) \ \Lambda_{ijk} \tag{8}$$

$$\lambda''_{ijk} = (4.4 \pm 0.4) \ \Lambda_{ijk},$$

where the errors correspond to the uncertainty in strong coupling constant: $\alpha_s(M_Z) = 0.12 \pm 0.01$. Inclusion of other uncertainties related $e.g.$ to threshold SUSY and GU corrections may require the doubling of the errors quoted. The renormalization effects due to third family Yukawa couplings [16] do not drastically change the relations (8). Let us define the renormalization factor $\eta/2$, relevant for proton decay as:

$$\lambda'(M_Z)\lambda''(M_Z) = \eta \cdot \Lambda^2. \tag{9}$$

From equation (8) we find: $\eta = 30 \pm 5$.

3 Proton decay due to R-parity violating couplings in the lowest order.

Simultaneous presence of both B- and L-violating couplings in GU models leads to proton decay. Let us consider the proton decay taking into account GU relations between couplings (6). There are two types of decay modes:

(1) $(B - L)$-conserving decays. The exchange of $\tilde{d}_i^c$ squarks between B-violating and L-violating vertices induces the 4-fermion operators:

$$2\eta \frac{\Lambda_{ijk}^* \Lambda_{ilm}}{m_{\tilde{d}_i}^2} \ \overline{d_j^c u_k^c} \ (V_{mn} d_n \nu_l - u_m e_l). \tag{10}$$

The kinematics selects the following 4-fermions operators in (10)

$$\overline{d^c u^c} \ \nu d, \quad \overline{d^c u^c} \ \nu s, \quad \overline{d^c u^c} \ eu, \quad \overline{d^c u^c} \ \mu u,$$

$$\overline{s^c u^c} \ \nu d, \quad \overline{s^c u^c} \ eu, \quad \overline{s^c u^c} \ \mu u \tag{11}$$

which lead to the proton decay. All these operators contain the u^c quarks, and therefore can be forbidden at tree level, if in the basis where u_i^c are the mass eigenstates we put

$$\Lambda_{ij1} = 0. \tag{12}$$

(2) $(B + L)$-conserving decays. The mixing of squarks: $\tilde{b}^c$, $\tilde{b}$ and $\tilde{t}^c$, $\tilde{t}$ leads to the operators:

$$2\eta \frac{(\Lambda_{3jk}\Lambda_{lmn}V_{n3})^*}{\mathcal{M}_{\tilde{b}}^2} \overline{d_j^c u_k^c} \, \overline{d_l^c \nu_m} \,,$$

$$-\eta \frac{(\Lambda_{ij3}\Lambda_{lm3})^*}{\mathcal{M}_{\tilde{t}}^2} \overline{d_i^c d_j^c} \, \overline{d_l^c e_m} \,,$$

(13)

where $\mathcal{M}_{\tilde{b}}^2$ and $\mathcal{M}_{\tilde{t}}^2$ parametrize the propagators $\tilde{b}^c - \tilde{b}$ and $\tilde{t}^c - \tilde{t}$ for low momenta. In particular,

$$\frac{1}{\mathcal{M}_{\tilde{b}}^2} = \frac{m_{\tilde{b}LR}^2}{m_{\tilde{b}LL}^2 m_{\tilde{b}RR}^2 - m_{\tilde{b}LR}^4},$$

(14)

where the mixing parameter $m_{\tilde{b}LR}^2$ is induced both via the μ-term at SUSY conserved level and via the soft breaking terms:

$$m_{\tilde{b}LR}^2 = m_b(A_b + \mu \tan\beta).$$

(15)

Here $A_b = O(m_{3/2})$ is soft breaking parameter. For $\tan\beta \sim 20 - 50$ the mixing mass may not be suppressed with respect to the diagonal masses $m_{\tilde{b}LL}^2$ and $m_{\tilde{b}RR}^2$. Consequently, the propagator factor $1/\mathcal{M}_{\tilde{b}}^2$ as well as $1/\mathcal{M}_{\tilde{t}}^2$ can be of the order of the factor $1/m_{\tilde{d}_i}^2$ from Eq. (10). We neglect the mixing of squarks from the lightest generations which are proportional to light quark masses. Mixing between squarks of different generation should be negligibly small to avoid the constraints from non-observation of flavour changing neutral-currents. Taking into account the kinematics we find from (13) the operators leading to proton decay:

$$\overline{d^c u^c} \, \overline{d^c \nu}, \quad \overline{s^c u^c} \, \overline{d^c \nu}, \quad \overline{d^c u^c} \, \overline{s^c \nu}, \quad \overline{d^c s^c} \, \overline{d^c \mu}.$$

The first three operators (with u^c) disappear if the conditions (12) are fulfilled; the last one can be removed by the equality

$$\Lambda_{123} = 0.$$

(16)

In fact, Λ_{ij1} and Λ_{123} may not be precisely zero; using relations (7) and renormalization effect (9) we get from (3) the bound on the GU scale couplings:

$$\Lambda_{ij1}, \quad \Lambda_{123} \lesssim 2 \cdot 10^{-13}.$$

(17)

In both conditions (12) and (16) the coupling constants with first family index are involved. Therefore we can assume the family hierarchy, according to which the couplings

6

with low indices (*i.e.* 1 and 2) are small, and maximal couplings are those with maximal number of family indices 3, first of all Λ_{233}, and then, probably, Λ_{133}. The question is: How large can be Λ_{233}?

4 Proton decay induced by Λ_{233} at one-loop.

Let us consider the configuration being the most protected from the proton decay, when there is only one term, $\Lambda_{233}\bar{5}_2\bar{5}_310_3$, in (4), with the following fermionic content of the supermultiplets: 10_3, includes t^c, $q_3 = (t, b')$, where $b' \equiv V_{3i}d_i$, and τ^c; $\bar{5}_3$ contains b^c and $l_3 = (\nu_\tau, \tau)$, $\bar{5}_2$ contains s^c and $l_2 = (\nu_\mu, \mu)$. All other terms in (4) have zero or negligibly small couplings. We will show that even in this case the proton decay appears as one-loop effect, thus leading to still very strong bounds on Λ_{233}.

Note that there is only one B-violating vertex: $\Lambda_{233}D_2^cD_3^cU_3^c$. It can be connected to L-violating vertex (needed to proton decay) by exchange of b^c, s^c or t^c and corresponding squarks. This allows to systematically find all relevant diagrams for proton decay. In accordance with (10) and (13) we get at the tree level the following $(B - L)$-conserving

$$2\eta \left|\Lambda_{233}\right|^2 \left[\frac{1}{m_{\tilde{s}}^2}\overline{b^ct^c}\left(b'\nu_\tau - t\tau\right) + \frac{1}{m_{\tilde{b}}^2}\overline{s^ct^c}\left(b'\nu_\tau - t\mu\right) \right] \tag{18}$$

and the $(B + L)$-conserving:

$$2\eta \left(\Lambda_{233}^*\right)^2 \left[\frac{V_{tb}^*}{\mathcal{M}_{\tilde{b}}^2}\overline{s^ct^c}\left(\overline{b^c\nu_\mu} - \overline{s^c\nu_\tau}\right) + \frac{1}{\mathcal{M}_{\tilde{t}}^2}\overline{s^cb^c}\left(\overline{b^c\mu} - \overline{s^c\tau}\right) \right] \tag{19}$$

operators. Also the operators are generated which can be obtained from (18) and (19) by replacement of two ordinary particles by their superpartners. The terms with $\tilde{s}^c - \tilde{s}$ exchange are omitted in (19), since they are proportional to small factor $V_{ts}/\mathcal{M}_{\tilde{s}}^2 \approx 10^{-3}/\mathcal{M}_{\tilde{b}}^2$ (in this equality we took into account that $\tilde{s}^c - \tilde{s}$ mixing is suppressed with respect to $\tilde{b}^c - \tilde{b}$ mixing by m_s/m_b, see (15), and therefore $\mathcal{M}_{\tilde{b}}^2/\mathcal{M}_{\tilde{s}}^2 \sim m_s/m_b$).

As we discussed before for kinematical reasons the operators (18) and (19) do not lead to proton decay. However, an additional exchange by the W-boson (or wino) as well as by charged Higgs (or Higgsino) converts the operators (18) and (19) (or the operators with superpartners) to the operators with light fermions which give proton decay already at one-loop level. Indeed, due to the presence of the CKM mixing the W- (wino), charged

Higgses (Higgsino) have family off-diagonal couplings (see Eq. (1)). The emission or absorbtion of these particles can reduce the generation index.

Let us find, using the operators (18) and (19), the crucial factors which appear in such a generation reduction:

(i) Evidently, the second term of (18) with four heavy fermions, and the fourth term with two t quarks, can not be transformed at one-loop into the operators with light particles only. Similarly, the third and the fourth terms in (19) stipulated by $\tilde{t}^c - \tilde{t}$ exchange do not give p-decay at one-loop. The third term contains two b^c quarks, the fourth one has three heavy fermions ($m > m_p$).

(ii). All the rest operators include t^c ($\tilde{t}^c$). The $t^c \to d$ conversion due to emission of charged Higgs boson or W-boson gives the factor $V_{td} \, m_t$ (in the case of the Higgs this factor follows from the Yukawa coupling (1); in the case of the W-exchange it comes from the chirality flip: $t^c \to t \to d \, W$). The same factor appears for $\tilde{t}^c \to d$ transition. Similarly, the conversion $t^c \to s$ ($\tilde{t}^c \to s$) implies the factor $V_{ts} m_t$.

(iii). The amplitudes of transitions of down quarks (squarks) $b^c \to u$ ($\tilde{b}^c \to u$), and $s^c \to u$ ($\tilde{s}^c \to u$) are proportional, respectively, to $V_{ub}^* \, m_b$ and $V_{us}^* \, m_s$. These factors are of the same order of magnitude.

(iv). L-violating part of $(B-L)$-operators (18) contains small CKM-elements V_{ts} or V_{td}, whereas $(B+L)$ operators (19) are proportional to $V_{tb} \approx 1$.

Combining the factors discussed in (ii)-(iv) we find that the largest one-loop amplitudes of $(B-L)$- and $(B+L)$-conserving modes contain the additional loop factors

$$
\begin{aligned}
\xi_{B-L} &= \frac{m_b m_t}{16 \pi^2 v^2} V_{ub}^* V_{td} V_{ts} \\
\xi_{B+L} &= \frac{m_b m_t}{16 \pi^2 v^2} V_{ub}^* V_{td} V_{tb}^*,
\end{aligned}
\tag{20}
$$

where $v \equiv \sqrt{v_1^2 + v_2^2}$, and $1/16\pi^2$ comes from loop integration. There are also transitions for which the loop factors can be obtained from (20) by substitution $V_{ub} \, m_b \to V_{us} \, m_s$.

Let us estimate the contributions from the leading diagrams.

1. The propagation of squark $\tilde{b}^c$ between B- and L-violating vertices (first term in (18)) "dressed" by charge Higgs (Higgsino) interaction leads to the diagrams shown in Fig. 1a,b. The mixing of charged Higgses H_1^-, H_2^+ in diagram of Fig. 1a and the coupling of three squarks in Fig. 1b are induced by soft SUSY breaking terms $\mu B H_1 H_2$

and $A\Lambda_{233}\tilde{s}^c\tilde{b}^c\tilde{t}^c$, correspondingly, where $A, B = O(m_{3/2})$. The estimation of diagrams gives

$$2\eta\frac{|\Lambda_{233}|^2}{m_{\tilde{b}}^2}\xi_{B-L}\left[du\ s\nu_\mu + su\ d\nu_\mu\right]. \tag{21}$$

In (21) we have taken into account the relation

$$\frac{\mu A}{v_1 v_2 m_{H+}^2} \sim \frac{\mu B}{v_1 v_2 m_{H+}^2} = \frac{1}{v^2}, \tag{22}$$

that connects the mass of the physical charged Higgs, m_{H+}^2, and the parameter of the mixing of scalar doublets H_1 and H_2, μB. The diagrams with dressing by W and wino (Fig. 1c,d) give similar result.

2. There are also the box diagrams with $\tilde{b}^c$ exchange, when H^+ emitted by t^c is absorbed by quark b' from L-violating vertex. Since $b' \to u^c$ transition is forbidden (b' couples to t^c, or t) the diagram gets the GIM suppression factor $V_{ub}m_b^2/\tilde{m}^2$, where $\tilde{m}^2$ is typical mass of squark. As the result one gets

$$\frac{A_{box}}{A_{vertex}} \propto \frac{m_u}{V_{ts}m_b}\frac{m_b^2}{\tilde{m}^2} < 10^{-4}.$$

Box diagrams lead to $(V - A)$-Lorentz structure of the effective operators.

3. The exchange of $\tilde{s}^c$ squark gives the diagram similar to those in Fig. 1a-d with emission of $\bar{\nu}_\tau$ instead of $\bar{\nu}_\mu$. The amplitude can be obtained from (21) by substitution

$$\frac{m_b V_{ub}^*}{m_{\tilde{s}}^2} \to \frac{m_s V_{us}^*}{m_{\tilde{b}}^2}. \tag{23}$$

4. "Dressing" the $(B + L)$ diagram with $\tilde{b}^c - \tilde{b}$ exchange (first and second terms in (19)) by $H^\pm$ and $\tilde{H}^\pm$ one gets the diagrams shown in Fig. 2a-d. Similar diagrams exist with dressing by W-boson and wino. The amplitudes corresponding to Fig. 2a,b can be estimated as:

$$2\eta\frac{(\Lambda_{233}^*)^2}{\mathcal{M}_{\tilde{b}}^2}\xi_{B+L}du\ \overline{s^c\nu_\tau}. \tag{24}$$

5. Box diagrams shown in Fig. 2c,d and similar diagrams with W and $\tilde{W}$ give the amplitudes comparable with that in (24) but having $(V - A)$-structure:

$$2\eta\frac{(\Lambda_{233}^2)^*}{\mathcal{M}_{\tilde{b}}^2}\xi_{B+L}\bar{s}^c\bar{\sigma}^\alpha d\ \bar{\nu}_\mu\bar{\sigma}_\alpha u. \tag{25}$$

6. The contributions of diagrams with exchange of $\tilde{s}^c - \tilde{s}$ (similar to those in Fig. 2a,b) are suppressed by factor of $m_s/m_b V_{ts}$, as we marked before.

According to previous discussion, the ratio of $(B + L)$- and $(B - L)$-amplitudes is

$$\frac{A_{B+L}}{A_{B-L}} \sim \frac{1}{V_{ts}} \frac{m_{\tilde{b}LR}^2}{m_{\tilde{b}}^2}. \tag{26}$$

For large $\tan\beta$ one has $m_{\tilde{b}LR}^2 \sim m_{\tilde{b}}^2$, and therefore $(B + L)$ amplitudes are enhanced by factor $1/V_{ts}$. Consequently, in models with Λ_{233} being the main source of R-parity violation the decay channels $p \to K^+\nu_\tau$ and $p \to K^+\nu_\mu$ dominate over $p \to \pi^+\overline{\nu_\tau}$, $p \to K^+\overline{\nu_\tau}$ channels (and similar modes with ν_μ). The $(B - L)$ channels may have branching ratios as small as $|V_{ts}|^2 \sim 10^{-3}$. In the case of $\tan\beta \sim 1$ the $(B - L)$- and $(B + L)$-amplitudes can be of the same order of magnitude.

Thus proton decay forbidden in the lowest order is generated due to CKM-mixing in one-loop. As follows from (21), (24) and (25) an additional suppression factor (20) appears in one loop amplitudes in comparison with tree level ones. Numerically it equals

$$\xi_{B+L} \equiv \xi = 5 \cdot 10^{-9} \left(\frac{m_b}{4.6 GeV}\right) \left(\frac{m_t}{176 GeV}\right) \left(\frac{V_{ub}^*}{3 \cdot 10^{-3}}\right) \left(\frac{V_{td}}{10^{-2}}\right). \tag{27}$$

Consequently, the bound on Λ_{233} can be relaxed by factor $\sqrt{\xi} \approx 7 \cdot 10^{-5}$:

$$\Lambda_{233} \lesssim 3 \cdot 10^{-9} \tag{28}$$

(compare with (17)). Using the amplitude (24) which can dominate at large $\tan\beta$ we find:

$$\Lambda_{233}^2 \lesssim 8 \cdot 10^{-18} \left(\frac{10^{-2}}{V_{td}}\right) \left(\frac{\mathcal{M}_{\tilde{b}}^2}{1 TeV^2}\right). \tag{29}$$

This result coincides with (28) at $V_{td} \sim 10^{-2}$ and $\mathcal{M}_{\tilde{b}}^2 \sim 1$ TeV.

Thus bouns on the proton lifetime strongly restrict even the Λ_{233} coupling of highest generations of matter fields. Large R-parity violating coupling constants are not admitted for any generation.

The following remarks are in order.

1). V_{td} is a common coefficient of all the amplitudes. For $V_{td} = 0$ one might have the suppression of all the one-loop contributions. However, the unitarity constraints of the CKM matrix gives for $V_{td} = (1 \pm 0.5) \cdot 10^{-2}$ at 90% C.L.

2). Lorentz structure of the one-loop operators differs from that of tree level operators. In particular, the vertex diagrams result in change of chirality (from right to left) of quarks from B-violating couplings. In box diagrams there is a change of chirality of one quark

from B-violating and one quark from L-violating couplings. Therefore no cancellation between one-loop and tree level contributions is expected.

3). The explicit computations of the diagrams confirm the results (21), (24) and (25), up to the factor

$$\frac{\ln x}{x - 1},\tag{30}$$

where $x \equiv m_t^2/m_{H+}^2$ for Higgs dressing, and $x \equiv m_t^2/m_W^2$ for W-dressing. W-contributions have also an additional factor 3. For $m_{H+} > 250$ GeV the contributions from Higgs dressed diagrams exceed those from diagrams with W.

4). Due to the relation (22) there is no dependence of amplitudes on μB, m_{H+}^2 or $\tan \beta$. This result is confirmed by explicit computation of diagrams up to the above factor (30).

5). Considerations performed in this section for Λ_{233} are valid for all couplings Λ_{ijk} which do not result in p-decay in the lowest order. For other couplings the bounds are even stronger. Thus (28) can be considered as the conservative bound on all R-parity violating coupling.

6). The analysis performed above and the bounds on R-parity violating constants are valid in more general context without grand unification. In (27) Λ_{233}^2 should be substituted by the product $|\lambda'_{233}\lambda''_{233}|$. Taking into account the renormalization effects we get at the electroweak scale:

$$|\lambda'_{233}\lambda''_{233}| \lesssim 5 \cdot 10^{-17} \left(\frac{10^{-2}}{V_{td}}\right) \left(\frac{\mathcal{M}_{\tilde{b}}^2}{1TeV^2}\right).\tag{31}$$

Similar or even stronger bounds can be obtained for any such products with at least two equal indices: $e.g.$ $|\lambda'_{ijk}\lambda''_{kjn}|$.

7). The presence of matter-Higgs mixing terms (4) does not change the bounds (28) unless strong fine tuning is implied.

5   Can R-parity violating couplings be large?

Let us consider the possibility to avoid the bound (28).

In the case of complex Higgs sector ($e.g.$ with additional 45-plets) there is a possibility to make another arrangement of particles in the $SU(5)$ multiplets. (In fact, such a sector allows to reproduce correct mass ratios m_e/m_d, m_μ/m_s). In principle, an arbitrary mixing (permutations) of the $SU(2) \times U(1)$ blocks from 5-plets as well as 10-plets of different

generates are admitted. In particular, in $\bar{5}_3$, together with b^c-quark it is possible to put some combination of the leptonic doublets:

$$L_3 \rightarrow (\hat{U}L)_3 \equiv \hat{U}_{3i}L_i, \tag{32}$$

and together with t^c in 10_3 one can put some combination of quark doublets:

$$Q_3 \rightarrow (\hat{W}Q)_3 \equiv \hat{W}_{3j}Q_j, \tag{33}$$

where $\hat{U}$ and $\hat{W}$ are arbitrary unitary matrices. Such a mixing of the $SU(2) \times U(1)$ blocks changes the structure of R-parity violating couplings, modifying the relation (7). In particular, for λ' and λ'' we get

$$\lambda'_{ijk} = 2\lambda''_{ij'k'}\hat{U}_{j'j}\hat{W}_{k'k} \tag{34}$$

instead of (7).

Let us find to what extend the rotations (32) and (33) can relax the bound on Λ. Suppose again that only λ''_{233} is non-zero. Dressing of the vertex $\lambda''_{233}(\tilde{s}^c b^c t^c - \tilde{b}^c s^c t^c)$ by Higgs (Higgsino) gives in one loop:

$$\lambda''_{233}\xi \left(\tilde{s}^c - k\tilde{b}^c \right) \left(\frac{V_{td}}{V_{ts}}ud + us \right), \tag{35}$$

where k is a constant of the order 1, and ξ is the one loop suppression factor (20). The coupling (35) does not depend on $\hat{W}$ and $\hat{U}$. The $\hat{W}$ and $\hat{U}$ rotations influence, however, the L-violating vertices. At tree level they become

$$\lambda''_{233}[S^c(\hat{W}D')_3(\hat{U}\nu)_3 - S^c(\hat{W}U)_3(\hat{U}E)_3 - B^c(\hat{W}D')_3(\hat{U}\nu)_2 + B^c(\hat{W}U)_3(\hat{U}E)_2]. \tag{36}$$

As we discussed before, in the case of $(B + L)$ conserving modes the squark $\tilde{s}^c$ emitted from the B-violating vertex mixes with $\tilde{s}$, and the latter is absorbed in the L-violating vertex (similarly, for $\tilde{b}^c$). According to (36) the amplitudes of the absorbtion of $\tilde{s}$ and $\tilde{b}$ are proportional to $\langle s|(\hat{W}d')_3\rangle$ and $\langle b|(\hat{W}d')_3\rangle$ respectively. Thus choosing $\hat{W}$ in such a way that $(\hat{W}d')_3 = d$ and suggesting that there is no flavour squark mixing (e.g. $\tilde{s}^c$ and $\tilde{d}$) one can suppress all $(B + L)$ decay modes in one-loop. Similar consideration holds for box diagrams.

The propagation of $\tilde{s}^c$ as well as $\tilde{b}^c$ between the vertices (35) and (36) results in the following $(B - L)$-conserving operators:

$$ud \left[(\hat{W}d')_3(\hat{U}\nu)_3 - (\hat{W}u)_3(\hat{U}e)_3 \right], \quad ud \left[-(\hat{W}d')_3(\hat{U}\nu)_2 + (\hat{W}u)_3(\hat{U}e)_2 \right]. \tag{37}$$

Since the neutrinos are massless (or very light) the only possibility to suppress the neutrino modes is to take $(\hat{W}d')_3 \equiv b$. Evidently, in this case the $(B + L)$-conserving modes are unsuppressed. Moreover, the equality $(\hat{W}d')_3 \equiv b$ means that $\langle(\hat{W}u)_3|u\rangle \equiv V_{ub}$ and the latter is non-zero. Consequently, second and fourth terms in (37) are not removed. Either $(\hat{U}e)_2$ or $(\hat{U}e)_3$ have an admixture of e or μ, and from (37) one gets, for instance, the operator $ud\ u\mu$ which leads to the proton decay.

Thus the additional rotations $\hat{W}$ and U do not allow to remove $(B - L)$ modes completely, but they change branching ratios, suppressing, $e.g.$, the neutrino modes. Eliminating the leading $(B + L)$ modes the $\hat{W}$ and $\hat{U}$ rotations relax the bound on Λ_{233} by factor $V_{ts}^{-1/2} \sim 5$.

Since CKM-mixing breaks any family symmetry, it is impossible to suppress the proton decay completely in the high orders of the perturbation theory. No horizontal symmetry can be introduced to forbid the operators of the type $(\bar{u}^c\bar{d}^c)(d\nu)$. B- and L-violation at least in some sector of the model will be propagated due to CKM-mixing to operators with light fermions which induce proton decay.

There are two evident possibilities to suppress proton decay:
1. suppress the mixing between matter generations;
2. modify the relation (7) between B- and L-violating couplings of usual matter fields in such a way that either B- or L- violating couplings are strongly suppressed (B-, L-violation asymmetry).

In the first case (since the mixing of known fermions is determined) one should introduce fourth fermion family, $\bar{5}_4$, 10_4 that has very small mixing with other families. For instance, the R-parity violating coupling $\Lambda_{234}\bar{5}_2\bar{5}_310_4$ generates the neutrino mass but does not result in fast proton decay, if the mixing with other generations is smaller than $V_{uB'} \sim 10^{-8}$. Note that B-violating coupling from the above term involves the quark of the fourth generation: $b^c s^c T^c$. To get the B-violating coupling $b^c s^c t^c$ without proton decay one should permute the fermions in multiplets in such a way that in 10_4 the upper quark T^c is substituted by t^c and in $\bar{5}_3$ the lepton doublet l_3 is substituted by l_4.

Concerning the second possibility, let us note that in grand unified theory with quark and lepton unification it is nontrivial to get the B- and L-violation asymmetry. As we will see in the next section the asymmetry can be related to the doublet-triplet splitting.

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6 R-parity violation and doublet-triplet splitting.

There are two possible ways to relate the asymmetry between the L- and B-violating couplings of usual matter fields in GU theories with doublet-triplet mass splitting.

1. Due to mixing of the matter and Higgs 5-plets (second term in (4)) the doublet-triplet splitting of the Higgs multiplet can lead to doublet-triplet asymmetry of matter field multiplet. This in turn breaks symmetry between quarks and leptons, and eventually, between the L- and B-violating couplings. Such a situation is realized in the model by Hall and Suzuki [23].

Let us consider an example of the model, where matter-Higgs mixing is the only source of R-parity violation. Suggesting as before that the third generation coupling dominates, we can write the appropriate terms of the superpotential in the following way

$$\bar{5}_3 \hat{m} H + \bar{H} \hat{M} H + \frac{m_{D,i}}{v_1} \bar{5}_i 10_i \bar{H}, \tag{38}$$

where $\bar{5}_i$ and 10_i are defined in the diagonal basis for down quark Yukawa couplings, so that d_i^c and d_i coincide, up to corrections M_W/M_{GU}, with mass eigenstates. The mass matrices of (38) can be written in the doublet-triplet form as:

$$\hat{m} = diag(m_{tripl}, m_{doubl}), \quad \hat{M} = diag(M_{tripl}, M_{doubl}), \tag{39}$$

where $M_{tripl} \sim M_{GU}$ and M_{doubl}, m_{doubl} and m_{tripl} are at the electroweak scale (large value of m_{tripl} would result in the fast proton decay). The first two terms in (38) can be removed by rotations of the doublet and the triplet components of the 5-plets $\bar{5}_3 = (B^c, L_3)$ and $\bar{H} = (\bar{T}, H_1)$. For triplet components we get the mixing:

$$\begin{aligned} \bar{T}' &= c_{tripl} \bar{T} + s_{tripl} B^c \\ B^{c\prime} &= c_{tripl} B^c - s_{tripl} \bar{T}, \end{aligned} \tag{40}$$

where $B^{c\prime}$ and $\bar{T}'$ are the mass states, $c_{tripl} \equiv \cos \alpha_{tripl}$, $s_{tripl} \equiv \sin \alpha_{tripl}$, and

$$\frac{s_{tripl}}{c_{tripl}} = \frac{m_{tripl}}{M_{tripl}}. \tag{41}$$

For doublet components:

$$\begin{aligned} H_1' &= c_{doubl} H_1 + s_{doubl} L_3 \\ L_3' &= c_{doubl} L_3 - s_{doubl} H_1, \end{aligned} \tag{42}$$

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and

$$\frac{s_{doubl}}{c_{doubl}} = \frac{m_{doubl}}{M_{doubl}}. \tag{43}$$

Since $m_{doubl}, m_{tripl}, M_{doubl} \sim M_W$ one gets from (43) and (41) that $s_{tripl} \sim M_W/M_{GU} < 10^{-14}$ is strongly suppressed, whereas s_{doubl} can be of the order 1.

Substituting the expressions (40) and (42) into the Yukawa coupling of (38) we obtain the effective R-parity violating couplings (2). In particular the third generation Yukawa coupling gives

$$\lambda_{333}^{eff} L_3 B^c Q_3', \tag{44}$$

where

$$\lambda_{333}^{eff} = s_{doubl} \frac{m_b}{v_1} = \frac{m_b}{v_1} \frac{m_{doubl}}{\mu}, \tag{45}$$

and $Q_3' \equiv V_{ib}^* Q_i$. Baryon violating interactions as well as pure leptonic terms are absent due to the antisymmetry. The Yukawa coupling of the second generation leads to

$$\frac{m_s}{v_1} V_{is}^* [s_{tripl} B^c S^c U_i^c + s_{doubl} L_3 S^c Q_i + s_{doubl} L_2 L_3 E_i^c] \tag{46}$$

(The first generation Yukawa coupling gives similar terms with substitution $m_s V_{is} \rightarrow m_d V_{id}$, $S \rightarrow D$, $L_2 \rightarrow L_1$).

The leading contribution to the proton decay is induced by L-violating interaction (44) and the B-violating coupling from (46). The $\tilde{b}^c$ exchange dressed by H^+, $\tilde{H}^+$... results in the amplitude for proton decay

$$A \propto \lambda_{333}^{eff} \cdot \frac{m_s}{v_1^2} s_{tripl} \; \xi_{B-L} = \frac{m_b m_s}{v_1^2} s_{doubl} s_{tripl} \; \xi_{B-L}. \tag{47}$$

Substituting values of parameters, we find that even for large $\tan \beta$ ($v_1 \sim m_b$) this amplitude is small enough to allow for s_{doubl}, and consequently, λ_{333}^{eff} to be of the order 1. All other diagrams give smaller contributions. (Note that in the considered example all the B-violating interactions contain b^c quark, so that even lowest family couplings need a loop "dressing").

Thus if the R-parity violation originates from the matter-Higgs mixing, the lepton violating coupling of the third generation can be large: $\lambda_{333}^{eff} = O(1)$.

The interaction (44) generates the neutrino mass in one loop:

$$m_{\nu_\tau} \sim \frac{|\lambda_{333}^{eff}|^2}{16\pi^2} \frac{m_b^2}{m_{\tilde{b}}^2} m_{3/2} = 0.1 MeV |\lambda_{333}^{eff}|^2 \left(\frac{m_{3/2}}{1TeV}\right) \left(\frac{1TeV}{m_{\tilde{b}}}\right)^2, \tag{48}$$

and for relatively light squarks the mass m_{ν_τ} can be easily in the MeV region. The cosmological bound on stable ν_τ : $m_{\nu_\tau} < 30$ eV gives the restriction $\lambda_{333}^{eff} < 10^{-2}$.

2. Another possibility to get the asymmetry of the B- and L-violation is to introduce the explicit doublet-triplet splitting in the matter multiplets. For this one should assume the existence of new superheavy matter fields.

Suppose that each generation of matter field contains an additional pair of 5-plets: $5'$ and $\bar{5}'$ with doublet-triplet splitting. For the third generation we introduce:

$$\bar{5}_3 = \begin{pmatrix} B^c \\ L_{3G} \end{pmatrix} \qquad \bar{5}_3' = \begin{pmatrix} B_G^c \\ L_3 \end{pmatrix} \qquad 5_3 = \begin{pmatrix} B_G \\ L_{3G}^c \end{pmatrix}, \tag{49}$$

where B_G^c, B_G, L_{3G}^c, and L_{3G} are new superheavy fields with mass $\sim M_{GU}$.

Note that by (49) we *generalize the doublet-triplet splitting* which is present now not only in the Higgs multiplets but also in the matter multiplets [1]. This "universal" doublet-triplet splitting could have an unique origin.

The electroweak symmetry breaking via the interaction $\bar{5}_3' 10_3 H$, results in mixing of the heavy and the light component with typical mixing angles:

$$\tan\alpha \sim \frac{M_W}{M_{GU}} \sim 10^{-14} \ . \tag{50}$$

Using the multiplets (49) one can introduce R-parity violating interactions even within one generation:

$$\Lambda_{333} \bar{5}_3 \bar{5}_3' 10_3. \tag{51}$$

This gives the terms:

$$B^c B_G^c T^c + B^c Q_3 L_3 - L_{3G} Q_3 B_G^c + L_{3G} L_3 \tau^c. \tag{52}$$

Note that there is no B-violating terms with only light matter fields. Mixing between B^c and B_G^c does not lead to such a term due to the antisymmetry of interaction. Proton decay

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[1]We will not discuss here the origin or the naturalness of the doublet-triplet splitting. Formally, the permutation of the light and the heavy matter fields can be achieved *e.g.* due to the interaction:

$$\bar{5}_3(M + h\Phi)5_3' + \bar{5}_3'(M' + h'\Phi)5_3,$$

where the parameters M, M', h, and h' are adjusted in such a way that B^c and L_3 are massless at the GU scale, whereas B_G^c, B_G, L_{3G}^c, and L_{3G} acquire masses $O(M_{GU})$.

is generated by one-loop diagram of the type shown in Fig. 1a with B^c being substituted by B^c_G. The corresponding suppression factor

$$\xi \frac{m_{\tilde{b}}^2}{M_{GU}^2} \ln \frac{m_H}{M_{GU}} \tag{53}$$

is strong enough to remove the bound on λ_{333}^{eff}. As in the previous case (44) the only R-parity violating coupling of light fields is the one with L-violation (λ'- type). It generates the neutrino mass (48).

Let us consider the possibility to get the B-violating coupling $s^c b^c t^c$. For this we introduce the additional 5-plets $\bar{5}'_2$ and $5'_2$ of second generation with new superheavy fermions S^c_G, S_G, L^c_{2G}, and L_{2G} and with permutation of light and heavy fermions, similar to that in (49). Now apart from the desired term $\bar{5}_2 \bar{5}_3 10_3$ one should admit also all other interactions which can be obtained from this by substitution $\bar{5}_2 \leftrightarrow \bar{5}'_2$ and $\bar{5}_3 \leftrightarrow \bar{5}'_3$:

$$(f_{333}\bar{5}_3\bar{5}'_3 + \lambda_{233}\bar{5}_2\bar{5}_3 + f_{233}\bar{5}'_2\bar{5}_3 + f_{323}\bar{5}_2\bar{5}'_3 + g_{233}\bar{5}'_2\bar{5}'_3 + f_{223}\bar{5}'_2\bar{5}_2)10_3. \tag{54}$$

(In fact, the permutation implies that the multiplets with permuted components have the same quantum numbers). However, if all these terms are present, they reproduce all the R-parity violating interactions (2) with light matter fields, and thus lead to the situation discussed in sect. 5. One possibility to solve the problem is to suggest strong hierarchy of couplings in (54). Also family symmetry can be introduced which forbids all the terms in (54) but the desired one. For instance, $U(1)$ symmetry with zero charge for $\bar{5}_2$, $\bar{5}_3$ 10_3 and charge 1 for all the rest multiplets make the desired selection. However, such a symmetry will be broken by mass terms, although this violation does not destroy the suppression of proton decay.

Let us finally remark that the doublet-triplet splitting breaks the $SU(4)$ symmetry responsible for $b - \tau$ unification at the GU scale. For instance in the model (38) the mass terms for bottom quark and tau lepton appear with the same couplings:

$$B^c Q_3 H_1 + \tau^c L_3 H_1, \tag{55}$$

but after the rotation (42) we get:

$$B^c Q_3 (c_{doubl} H'_1 - s_{doubl} L'_3) + \tau^c L'_3 H'_1. \tag{56}$$

Consequently the generation of the R-parity violating coupling $L_3 B^c Q_3$ with constant proportional to s_{doubl} turns out to be connected to the reduction of the $b - \tau$ mass ratio by the factor c_{doubl}.

7   Discussion and Conclusions.

1. The R-parity violating couplings may have strong flavour hierarchy, so that the coupling constants for the fields from the third generation could be of the order 1. These couplings may have a number of phenomenological consequences: generation of the MeV mass of ν_τ, change of the infrared fixed point of the top quark, renormalization of the mass ratio m_b/m_τ, etc..

2. Motivated by the success of the supersymmetric Grand Unification, we have considered the possibility of existence of such large couplings in the Grand Unified theories. In the lowest order of perturbation theory the bound from the proton decay can be satisfied by smallness or absence of couplings for low generations. However, being suppressed in the lowest order the proton decay appears inevitably at one-loop as the consequence of the CKM-mixing. In the safest case with only one nonzero coupling Λ_{233} the bound (28) $\Lambda_{233} \lesssim 3 \cdot 10^{-8}$ has been obtained, which can be considered as the conservative bound on all R-parity violating couplings in $SU(5)$ models.

3. The analysis and the bounds obtained here are valid in a more general context. They correspond to the bounds on the product of B- and L-violating couplings $\lambda'\lambda'' \lesssim 5 \cdot 10^{-17}$, where λ' and λ'' have at least one common index.

4. In models with R-parity violation, especially in the case of one-loop induced decay, the proton decay modes may differ from those in the usual supersymmetric model. In particular, the modes with $(B + L)$-conservation, like $p \rightarrow K^+ \nu_\mu$ and $n \rightarrow K^+ \mu^-$, can dominate over the $(B - L)$-conserving modes, like $p \rightarrow K^+ \nu_\tau$ and $p \rightarrow K^0 \mu^+$.

5. The bound (28) can be avoided if new fermions (new matter fields) exist which mix very weakly with known fermions. These could be the fermions from the fourth generation.

The bounds can also be avoided if there is an asymmetry of B- and L-violating interactions, namely if either L- or B-violating interactions are strongly suppressed. This asymmetry can be related to the doublet-triplet splitting. In the simple examples the largest R-parity violating coupling is the one with L-violation.

6. For coupling constants Λ satisfying the bound (28), no appreciable effects of R-parity violation in accelerator experiments are expected. Also, the generated neutrino masses are very small. Inversely, the observation of R-parity violating effects at accelerators will have strong impact on the Gran Unification: this can imply Higgs-matter mixing or doublet-triplet splitting in matter supermultiplets.

Acknowledgements.

Figure Captions

Fig. 1: Leading one-loop diagrams of $(B - L)$-conserving p-decay in the model with $\Lambda_{233} \neq 0$. Similar diagrams exist with $\tilde{s}^c$ exchange and the emission of $\bar{\nu}_\tau$.

Fig. 2: Leading one-loop diagrams of $(B + L)$-conserving p-decay in the model with $\Lambda_{233} \neq 0$. Similar diagrams exist with substitution $H \rightarrow W$, $\tilde{H} \rightarrow \tilde{W}$.

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