Published by Institute of Physics Publishing for SISSA/ISAS

Received: *April 6, 2002*
Accepted: *June 14, 2002*

U(1) flavor symmetry and proton decay in supersymmetric standard model

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Abstract: Issues of proton decay are examined in supersymmetric standard model with U(1) flavor symmetry. Dimension five proton-decay operators which arise generically are controlled by the flavor symmetry. We show that unlike the minimal supersymmetric SU(5) case the proton decay modes containing charged lepton can have large branching ratios if the dimension five operators of left-handed type dominate. Measuring the branching ratio of the electron mode to the muon mode may reveal the mechanism of the neutrino mass generation. The case with vanishing charges for Higgs doublets marginally survives the present experimental bound on proton life time.

Keywords: Supersymmetry Phenomenology, Quark Masses and SM Parameters, Supersymmetric Standard Model.

http://jhep.sissa.it/archive/papers/jhep062002032/jhep062002032.pdf

JHEP06(2002)032

Contents

1. Introduction

With particle contents and gauge symmetry of the standard model, nucleon decay may take place through non-renormalizable interaction. In fact, the gauge invariance allows dimension six operators with baryon and lepton number violation. They may be induced by exchange of X, Y gauge bosons and colored Higgs multiplets in a grand unified theory (GUT), or may just arise as non-renormalizable operators at a fundamental scale which we assume to be near the Planck scale. The operators of the latter origin are, however, suppressed by the Planck mass squared, and thus the nucleon decay induced by them is suppressed to a level where near future experiments will not be able to reach.

The situation is very different if one considers the supersymmetric (SUSY) extension of the Standard Model. Existence of superparticles allows gauge invariant dimension five operators which can induce nucleon decay after superparticle dressing [1].[1] These operators are very dangerous because they are suppressed only by a single power of the Planck mass. In fact, for the superparticle masses around 1 TeV, present proton decay experiments constrain the mass scale of the dimension five operators much larger than the Planck mass, or in other words, their coefficients should be much smaller than unity when normalized by the Planck scale. This is indeed embarrassing if one believes the widely accepted argument on the generality that all operators which are allowed by symmetry should arise with order one coefficients. There are many attempts to explain the smallness of these generic dimension 5 operators. They include 1) imposing some symmetry such as the family

[1] We assume R-parity conservation, and thus we do not consider the most dangerous dimension four operators.

symmetry [2]–[6], the discrete gauge symmetry [7], the Peccei-Quinn symmetry [8, 9], the $\mathrm{U}(1)_A$ symmetry [10] and the R symmetry [11, 12], and 2) attributing to configurations of quarks and leptons in extra spatial dimensions [13, 14].

In this paper, we shall re-examine the case of U(1) flavor symmetry [15] and closely investigate how the nucleon decays can be (or cannot be) suppressed in the context of the minimal supersymmetric standard model (MSSM). Furthermore, we will explore the correlation between the branching fractions and the family structure [2, 3, 4].

We will show that, unlike a typical SUSY SU(5) GUT, the proton decay modes containing charged lepton can have sizable branching ratios and the ratio of the μ mode and the e mode may give us a crucial information on the neutrino mass generation. We will also study the proton life time. Keeping in mind that our argument based on the flavor symmetry contains uncertainties in coefficients of operators, we will argue that the case with vanishing charges for the Higgs doublets marginally survives the present proton decay bound. The constraint is relaxed for negative Higgs charges. We will briefly discuss how to generate higgsino mass μ in this case.

2. U(1) Flavor symmetry and charge assignment

A hypothetical flavor symmetry is a symmetry to explain the Yukawa structure. Here we consider the Froggatt-Nielsen mechanism [15] in which we assign Froggatt-Nielsen U(1) charges q to the MSSM superfields Q and -1 to the Froggatt-Nielsen field X. The U(1) symmetry is spontaneously broken when the X field develops a vacuum expectation value. We define λ as $\lambda \equiv \langle X\rangle/M_{\mathrm{pl}}$, with M_{pl} being the Planck scale. In the MSSM, we have the following superpotential

$$W = y_U^{ij} Q_i U_j^C H_u + y_D^{ij} Q_i D_j^C H_d + y_E^{ij} E_i^C L_j H_d \,, \tag{2.1}$$

where $i, j = 1, 2, 3$ are generation indices. y_U^{ij} etc. are Yukawa couplings, whose magnitudes are governed by the U(1) symmetry

$$y_U^{ij} = f_U^{ij} \lambda^{q_i + u_j^c + h_u} \,, \qquad y_D^{ij} = f_D^{ij} \lambda^{q_i + d_j^c + h_d} \,, \qquad y_E^{ij} = f_E^{ij} \lambda^{e_i^c + l_j + h_d} \,. \tag{2.2}$$

Here f_U^{ij} etc. are somewhat arbitrary constants which are typically of order unity. In order to obtain physical masses and mixings, we must translate fields from a flavor basis into a mass basis through unitary matrices as

$$U_Q^{u\mathrm{T}} y_U U_U = \mathrm{diag}(y_u, y_c, y_t) \,, \tag{2.3}$$
$$U_Q^{d\mathrm{T}} y_D U_D = \mathrm{diag}(y_d, y_s, y_b) \,, \tag{2.4}$$
$$U_E^{\mathrm{T}} y_L U_L = \mathrm{diag}(y_e, y_\mu, y_\tau) \,, \tag{2.5}$$

from which we obtain the CKM matrix as $V_{\mathrm{KM}} = U_Q^{u\dagger} U_Q^d$.

Following the conventional wisdom, we identify λ with the Wolfenstein parameter $\lambda \sim 0.22$, and determine the U(1) charges of the MSSM fields. The experimental values of the masses and mixings near the Planck scale are approximated as [16]

$$V_{us} \sim \lambda \,, \qquad V_{cb} \sim \lambda^2 \,, \qquad V_{ub} \sim \lambda^3 \,, \tag{2.6}$$

and

$$
\begin{aligned}
m_u : m_c : m_t &\sim \lambda^8 : \lambda^4 : 1 \,, && (2.7)\\
m_d : m_s : m_b &\sim \lambda^4 : \lambda^2 : 1 \,, && (2.8)\\
m_e : m_\mu : m_\tau &\sim \lambda^5 : \lambda^2 : 1 \,. && (2.9)
\end{aligned}
$$

In this paper, we will consider the following class of charge assignments:

$$
\begin{aligned}
q_1 &= q_3 + 3 \,, \qquad q_2 = q_3 + 2 \,, && (2.10)\\
u_1^c &= u_3^c + p \,, \qquad u_2^c = u_3^c + 2 \,, && (2.11)\\
e_1^c &= e_3^c + r - n - m \,, \qquad e_2^c = e_3^c + 2 - n \,, && (2.12)\\
d_1^c &= d_3^c + 1 \,, \qquad d_2^c = d_3^c \,, && (2.13)\\
l_1 &= l_3 + n + m \,, \qquad l_2 = l_3 + n \,, && (2.14)
\end{aligned}
$$

where n, m, p and r are integers with $0 \leq n \leq 2$, $0 \leq m \leq 5-n$, $3 \leq p \leq 5$ and $4 \leq r \leq 5$. Here, eq. (2.10) is dictated by the magnitudes of the CKM matrix elements, eq. (2.11) and eq. (2.13) are obtained from the up- and down-type quark mass ratios and the mixings, and eq. (2.12) and eq. (2.14) are determined by the charged lepton mass hierarchy. Uncertainty of the coefficients f_U^{ij} etc. allows some ambiguities of the charges, which are parameterized by p and r. Furthermore it follows from the fact that $m_b \sim m_\tau$ and $y_t \sim 1$

$$
q_3 + d_3^c = e_3^c + l_3 \,, \qquad q_3 + u_3^c + h_u = 0 \,. \tag{2.15}
$$

The parameters m and n are determined by the oscillation of the atmospheric [17] and solar [18] neutrinos. Here we assume that the see-saw mechanism [19] operates to generate neutrino masses and mixings. The overall scale of the neutrino masses is fixed by the scale of lepton number violation, and the ratio of the masses as well as the mixing angles are controlled by the U(1) charges. We will consider the following two cases, both of which reproduce the bimaximal mixing angles:
i) the lopsided type [20]

$$
n = 0, \qquad m = 1 \,, \tag{2.16}
$$

and
ii) the anarchical type [21, 22]

$$
n = 0 \,, \qquad m = 0 \,, \tag{2.17}
$$

with the right-handed neutrinos having appropriate U(1) charges.

To make our argument explicit, we will consider the following two types of charge assignment.

1. *Best-fit* charge assignment. In the first type of charge assignment, we take

$$
p = 5 \,, \qquad r = 5 \,, \qquad n = 0 \,, \qquad m = 0 \text{ or } 1 \,, \tag{2.18}
$$

in eqs. (2.10)–(2.14). We call this assignment the *best-fit* one. Notice that it is not consistent with the GUT symmetry where the Q and U^C, for instance, are in a single multiplet and thus have a common U(1) charge.

2. *GUT-inspired* charge assignment. The second type of charge assignment we will discuss is the one consistent with the SU(5) GUT. Specifically we consider

$$q_i = u_i^c = e_i^c\,, \quad q_1 = q_3 + 3\,, \quad q_2 = q_3 + 2\,, \quad (p = 3, r = 4)\,, \tag{2.19}$$
$$d_i^c = l_i\,, \quad d_1^c = d_3^c + 1\,, \quad d_2^c = d_3^c \quad (n = 0, m = 1)\,. \tag{2.20}$$

The charge assignment for u^c would predict the up-quark Yukawa coupling of λ^6, much larger than the actual value. Thus in this scheme, the small up-quark mass should be attributed to either accidentally small coefficients f_U^{i1} of order λ^2 or the cancellation between the matrix elements.

3. Proton decay

Now we would like to discuss proton decay. Gauge invariance allows the dimension five operators in the superpotential, which break the baryon number as well as the lepton number conservation and are divided into an LLLL part and an RRRR part:

$$W = \frac{1}{2M} C_L^{ijkl} Q_i Q_j Q_k L_l + \frac{1}{M} C_R^{ijkl} E_i^C U_j^C U_k^C D_l^C\,. \tag{3.1}$$

Here C_L^{ijkl} and C_R^{ijkl} are some coefficients, and M denotes the fundamental scale of the theory, which we identify with the Planck scale $M_{\rm pl}$. If these coefficient were of order unity, they would cause proton decay with life time several magnitudes shorter than the present experimental bound. In the framework of the Froggatt-Nielsen U(1) symmetry, we expect suppression

$$C_L^{ijkl} \sim \lambda^{q_i+q_j+q_k+l_l} \left(\sim y_U^{ij} y_D^{kl} \lambda^{-h_u-h_d}\, \lambda^{q_j-u_j^c+l_l-d_l^c}\right), \tag{3.2}$$
$$C_R^{ijkl} \sim \lambda^{e_i^c+u_j^c+u_k^c+d_l^c} \left(\sim y_U^{ij} y_D^{kl} \lambda^{-h_u-h_d}\, \lambda^{e_i^c-q_i+u_k^c-q_k}\right). \tag{3.3}$$

Here we have omitted constants in front of the powers of λ. It is interesting to note here that the proton decay is suppressed if the sum of the two Higgs charges is negative, $h_u + h_d < 0$. We will come back to this point later on.

After wino and higgsino dressing,[2] we obtain proton decay amplitudes [24]–[29][3]

$$\text{Amp}(p \to K^+ \bar{\nu}_e) \sim \frac{f}{M} \left(A_1^1 + A_2^1\right), \tag{3.4}$$
$$\text{Amp}(p \to K^+ \bar{\nu}_\mu) \sim \frac{f}{M} \left(A_1^2 + A_2^2\right), \tag{3.5}$$
$$\text{Amp}(p \to K^+ \bar{\nu}_\tau) \sim \frac{f}{M} \left(A_1^3 + A_2^3 + A_3 + A_4\right), \tag{3.6}$$

[2]We assume that flavor mixings in the quark-squark-gluino interactions are small enough to avoid the dangerous flavor changing processes, and simply ignore gluino dressing diagrams in this paper. If these generation mixings are sizable, the gluino dressing diagrams give considerable contributions to proton decay [23].

[3]Proton decay to pion suffers from Cabbibo suppression, and thus we will not consider it in this paper.

where[4]

$$A_1^k \sim g_2^2 V_{cd} V_{cs} C_L^{221k}\,, \tag{3.7}$$

$$A_2^k \sim g_2^2 V_{td} V_{ts} C_L^{331k}\,, \tag{3.8}$$

$$A_3 \sim y_t y_\tau V_{ts} C_R^{3311}\,, \tag{3.9}$$

$$A_4 \sim y_t y_\tau V_{td} C_R^{3312}\,. \tag{3.10}$$

In eqs. (3.7) and (3.8), we have suppressed other contributions proportional to C_L^{112k}, C_L^{113k}, etc., which are of the same order as the ones explicitly written there. In eqs. (3.4)–(3.6) f represents a loop factor, $f \sim 1/(16\pi^2 m_{\mathrm{SUSY}})$, where m_{SUSY} is a representative sparticle mass scale. Degeneracy of the squark and slepton masses in different generations is implicitly assumed in the above evaluation. That is, we do not consider the so-called effective supersymmetry where the squarks and sleptons in the first two generations are very heavy.

The amplitudes of the decay $n \to K^0 \bar{\nu}$ are the same as those of $p \to K^+ \bar{\nu}$ up to factors of order unity.

Using a numerical relation

$$y_t y_\tau = \frac{m_t m_\tau}{2 m_W^2} g_2^2 \left(\tan\beta + \frac{1}{\tan\beta} \right) \sim \frac{1}{2} \lambda^2 g_2^2 \tan\beta, \quad \tan\beta \gtrsim 3 \tag{3.11}$$

at the Z-boson mass scale, we can rewrite the coefficients A as

$$A_1^k \sim A_2^k \sim g_2^2 y_t y_b \, \lambda^{-h_u - h_d + \delta + 8 + (l_k - l_3)}\,, \tag{3.12}$$

$$A_3 \sim A_4 \sim g_2^2 y_t y_b \, \lambda^{-h_u - h_d - \delta + 5 + p} \, \frac{\tan\beta}{2}\,, \tag{3.13}$$

where $\delta \equiv 2q_3 - u_3^c - e_3^c$. Notice that the δ dependence of the LLLL operators and that of the RRRR operators are inverse.

As for charged lepton modes, we obtain

$$\mathrm{Amp}(p \to K^0 e^+) \sim \frac{f}{M} A_5^1\,, \tag{3.14}$$

$$\mathrm{Amp}(p \to K^0 \mu^+) \sim \frac{f}{M} A_5^2\,, \tag{3.15}$$

where

$$A_5^k = g_2^2 \sum_i C_L^{11ik} V_{is} \sim g_2^2 y_t y_b \, \lambda^{-h_u - h_d + \delta + 8 + (l_k - l_3)}\,. \tag{3.16}$$

Since the proton cannot decay to the tau lepton, the decay amplitudes do not contain the third generation down-type Yukawa coupling constant. Therefore, contributions from the RRRR operators to the charged lepton modes are negligible.

[4] C_L^{ijkl} in eqs. (3.7) and (3.8) should be replaced by $C_L^{mnpq} (U_Q^u)_m{}^i (U_Q^u)_n{}^j (U_Q^u)_p{}^k (U_L)_q{}^l$. However, since the off-diagonal elements of the mixing matrices are suppressed by some powers of λ (including λ^0), we deduce that $C_L^{mnpq} (U_Q^u)_m{}^i (U_Q^u)_n{}^j (U_Q^u)_p{}^k (U_L)_q{}^l \sim C_L^{ijkl}$ in this framework. The same argument is applied to C_R.

3.1 Branching ratios

Let us first argue the branching ratios of the proton decay. When the LLLL operators dominate in the decay modes to the neutrinos, we find from the above consideration

$$\frac{\Gamma\left(p \to K^0 l_k^+\right)}{\Gamma(p \to K^+ \bar{\nu}_k)} \sim \mathrm{O}(1)\,, \qquad (k = 1, 2)\,. \tag{3.17}$$

Thus, we predict that

$$\frac{\Gamma\left(p \to K^0 e^+\right)}{\Gamma(p \to K^+ \bar{\nu})} \sim \lambda^{2(n+m)}\,, \qquad \frac{\Gamma\left(p \to K^0 \mu^+\right)}{\Gamma(p \to K^+ \bar{\nu})} \sim \lambda^{2n}\,, \tag{3.18}$$

where $\Gamma(p \to K^+\bar{\nu}) = \sum_k \Gamma(p \to K^+\bar{\nu}_k)$ is the decay width into three types of the neutrinos. In particular, for $n = 0$, which is strongly suggested by the large mixing between ν_τ and ν_μ, we conclude that the decay rate of the muon mode is comparable to that of the neutrino mode. Furthermore, the decay rate to the electron is controlled by the U(1) charge of the l_1, and thus we may be able to distinguish the anarchical type assignment for the leptons from the lopsided one by measuring the ratio $\Gamma(p \to K^0 e^+)/\Gamma(p \to K^0 \mu^+)$.

This consequence is a striking contrast to the case of the minimal SU(5) SUSY GUT in which the proton decay operators are induced by exchanging the colored Higgs multiplets. In fact, in the SU(5) case the charged lepton modes are negligibly suppressed [25, 26]. This is because the decay amplitudes into the charged leptons are proportional to the very small up-quark mass, while those into the neutrinos can be proportional to the charm- or top-quark mass.

On the other hand, when the RRRR operators are dominant, the proton mainly decays through $p \to K^+\bar{\nu}_\tau$. Thus we expect

$$\frac{\Gamma\left(p \to K^0 e^+, K^0 \mu^+\right)}{\Gamma(p \to K^+ \bar{\nu})} \ll \mathrm{O}(1)\,. \tag{3.19}$$

Which of the operators, LLLL or RRRR, dominates depends on the charges one assumes. When we take the *best-fit* charge assignment, $l_2 - l_3 = 0$ and $p = 5$, whereas δ is a parameter which is not fixed. Then, for

$$\delta \lesssim 0.9 - 0.3 \ln(\tan\beta/3)\,, \tag{3.20}$$

the LLLL contributions dominate over the RRRR ones.

In the GUT-inspired case, $l_2 - l_3 = 0$, $p = 3$ and $\delta = 0$.[5] It follows from eq. (3.13) that for small $\tan\beta$, the LLLL amplitudes are of the same order as the RRRR amplitudes. Thus, we expect that the muon mode is comparable to the neutrino mode and the electron

[5]If the theory is really embedded into a GUT group, one has to take into account the contribution from the colored Higgs exchange. In the GUT framework, our argument given here is valid when some mechanism operates to suppress the proton decays mediated by the colored Higgses and thus the genuine dimension five operators controlled by the flavor symmetry are the dominant sources of the proton decay. The suppression mechanisms of the proton decay operators coming from the exchange of the colored Higgses have been proposed, for instance, in refs. [9, 10, 30, 31]. See also ref. [32] and references therein.

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mode is suppressed by λ^2. Namely,

$$\frac{\Gamma\left(p \to K^0\mu^+\right)}{\Gamma(p \to K^+\bar{\nu})} \sim 1\,, \qquad \frac{\Gamma\left(p \to K^0 e^+\right)}{\Gamma(p \to K^+\bar{\nu})} \sim \lambda^2 \quad (\tan\beta \sim 3)\,. \tag{3.21}$$

On the contrary, for large $\tan\beta$, the RRRR contributions are enhanced and thus the neutrino mode, $p \to K^+\bar{\nu}$, dominates.

3.2 Decay rates

Next we will derive constraints on the charge assignment of the MSSM fields from the null results of the proton decay searches. In the framework we are considering, the decay rates for the modes $p \to K^+\bar{\nu}, K^0e^+, K^0\mu^+$ and $n \to K^0\bar{\nu}$ can be comparable in magnitudes. However, since the most severe experimental constraint comes from the decay mode $p \to K^+\bar{\nu}$, all we have to do is to study this mode.

The partial decay rate for the mode $p \to K^+\bar{\nu}$ is calculated as

$$\Gamma(p \to K^+\bar{\nu}) \sim \frac{m_p}{32\pi}\left(1 - \frac{m_{K^+}^2}{m_p^2}\right)^2 \frac{1}{f_\pi^2}\left(\sum_k \left|R_L\beta_p\,\mathrm{Amp}_{Lk} + R_R\alpha_p\,\mathrm{Amp}_R\,\delta_{k3}\right|^2\right), \tag{3.22}$$

where

$$\mathrm{Amp}_{Lk} \sim \frac{f}{M} g_2^2 y_t y_b\, \lambda^{-h_u - h_d + \delta + 8 + (l_k - l_3)}\,, \tag{3.23}$$

$$\mathrm{Amp}_R \sim \frac{f}{M} g_2^2 y_t y_b\, \lambda^{-h_u - h_d - \delta + 5 + p}\, \frac{\tan\beta}{2}\,. \tag{3.24}$$

Here, m_p and m_{K^+} are the masses of proton and K^+ respectively, f_π is the pion decay constant, $R_{L,R}$ represent renormalization effects of LLLL and RRRR operators from M to 1 GeV, and α_p and β_p are the hadronic matrix element parameters. Hereafter, for simplicity we fix the parameters as $m_{\mathrm{SUSY}} = 1\,\mathrm{TeV}$, $|\alpha_p| = 0.015\,\mathrm{GeV}^3$, $|\beta_p| = 0.014\,\mathrm{GeV}^3$ [33], $R_L = 10.2$ and $R_R = 6.5$.[6]

For the *best-fit* charge assignment, in order not to conflict with the experimental bound $\tau(p \to K^+\bar{\nu}) > 1.9 \times 10^{33}$ yr [34], the U(1) charges must satisfy the following inequalities:

$$y_t y_b\, \lambda^{-h_u - h_d + \delta + 8} \lesssim 4.2 \times 10^{-8}\,, \tag{3.25}$$

$$y_t y_b \lambda^{-h_u - h_d - \delta + 10}\, \frac{\tan\beta}{2} \lesssim 6.2 \times 10^{-8}\,. \tag{3.26}$$

Substituting the approximate values $m_t \sim 110\,\mathrm{GeV}$ and $m_b \sim 1.0\,\mathrm{GeV}$ near the Planck scale, the above are rewritten as

$$\lambda^{-h_u - h_d + \delta}\, \tan\beta \lesssim 2.1\,, \tag{3.27}$$

$$\lambda^{-h_u - h_d - \delta}\, \frac{\tan^2\beta}{2} \lesssim 64\,. \tag{3.28}$$

[6] $R_{L,R}$ are evaluated by solving one-loop renormalization group equations due to gauge interactions from a high energy scale ($\sim 2 \times 10^{16}$ GeV) to 1 GeV. Inclusion of the Yukawa couplings and effects above this high energy scale do not qualitatively change our results.

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Thus, we obtain

$$h_u + h_d \lesssim -0.2 + \delta - 0.7 \ln(\tan\beta/3)\,, \tag{3.29}$$

$$h_u + h_d \lesssim 1.8 - \delta - 1.3 \ln(\tan\beta/3)\,. \tag{3.30}$$

The proton decay constraint is relaxed most when $\tan\beta \sim 3$ and $\delta = 1$, which gives the upper bound

$$h_u + h_d \lesssim 1\,, \tag{3.31}$$

and hence we find that the case $h_u + h_d = 0$ survives the proton decay constraint for very low $\tan\beta$. On the other hand, for larger $\tan\beta$, the sum of the two Higgs charges must be negative for any choice of δ to satisfy the proton decay bound.

For the GUT-inspired case where $\delta = 0$ and $p = 3$, the constraints on the Higgs charges are modified as

$$h_u + h_d \lesssim -0.2 - 0.7 \ln(\tan\beta/3)\,, \tag{3.32}$$

$$h_u + h_d \lesssim -0.3 - 1.3 \ln(\tan\beta/3)\,. \tag{3.33}$$

Enhancement of the RRRR contributions arises from the fact that difference between u_1^c and u_3^c is smaller than that in the *best-fit* case. Thus we conclude that the charge assignment $h_u + h_d = 0$ is marginally allowed for very small $\tan\beta$.

In deriving the constraints above, one has to keep in mind that there are some uncertainties in evaluating the proton decay rate. They include

1. The coefficients of the proton-decay dimension five operators which have been set order unity may be accidentally small, like m_u in the GUT-inspired case.
2. There are some uncertainties in α_p and β_p. Taking the smallest allowed values for them, the proton decay amplitudes become five times smaller than our estimate.
3. We have taken the representative sparticle mass $m_{SUSY} = 1\,\text{TeV}$. However, it becomes larger and the proton decay gets suppressed if we adopt smaller wino and higgsino masses and larger squark and slepton masses.

These could reduce the proton decay rate, and then the constraints for the Higgs charges obtained above would become somewhat relaxed.

On the other hand, negative charge assignment for the Higgs doublets [30] survives the proton decay constraint in a wider region of the parameter space as the proton decay rate is reduced by $\lambda^{-2(h_u+h_d)}$.

4. μ parameter

Here we would like to make a brief comment on the higgsino mixing parameter μ when h_u+h_d is negative. In this case, neither the term $[\mu H_u H_d]_{\theta^2}$ nor $[\mu'(X/M_{\rm pl})^n H_u H_d]_{\theta^2}$ (with $n > 0$) is invariant under the Froggatt-Nielsen U(1) symmetry. Thus the generation of the μ term is somewhat contrived. Consider the Giudice-Masiero mechanism [35] with the term $[Z^\dagger/M_{\rm pl}(X^\dagger/M_{\rm pl})^{-(h_u+h_d)} H_u H_d]_{\theta^2\bar\theta^2}$. This gives the μ term of $F_Z^\dagger/M_{\rm pl}\lambda^{-(h_u+h_d)} \sim m_{3/2}\lambda^{|h_u+h_d|}$. This is much smaller than the electroweak scale unless the gravitino mass

$m_{3/2}$ is large. This is, in fact, the case in, e.g. anomaly mediation where $m_{3/2}$ is around a few tens TeV. Thus as far as $\lambda^{|h_u+h_d|} \gtrsim 10^{-2}$, we obtain the weak scale μ parameter. However, the Higgs mixing parameter, B, is not suppressed and in general as heavy as $m_{3/2}$. Then we will have difficulty in obtaining the correct electroweak symmetry breaking.

A possible mechanism is to introduce a singlet S with a positive U(1) charge so that $[SH_uH_d]_{\theta^2}$ is invariant under the flavor symmetry. It is assumed that S field does not have a VEV in the SUSY limit. SUSY breaking generally generates a VEV of S, which is nothing but the μ [36]. An explicit realization was given in ref. [37], where the charges of relevant fields should be adjusted to obtain the μ parameter of correct order of magnitude. Thus the choice of a negative $h_u + h_d$ does not immediately cause trouble to generate the μ parameter, though the way to do is rather restricted.

5. Conclusions and discussion

In this paper, we have considered the proton decay when the dimension five operators are controlled by the U(1) flavor symmetry. We have shown that the charged lepton decay modes can have sizable branching ratios when the LLLL operators dominate over the RRRR ones. In this case, measurement of the branching ratio of the e mode and the μ mode will provide an important information on the charges for the doublet leptons and thus reveal how the neutrino masses are generated. On the other hand, if the RRRR operators are the main sources of the proton decay, the neutrino mode dominates in the proton decay. We have also investigated the proton decay rate and showed that the proton-decay suppression due to the U(1) flavor symmetry marginally survives the present proton decay bound even when the Higgs doublets do not have the U(1) charges. The bound is less severe, if we take negative charges of the Higgs doublets.

We have made several assumptions in deriving the conclusions drawn above. Firstly we have assumed that non-renormalizable operators are suppressed by the four-dimensional Planck scale. However, there are many models in which the fundamental scale is lower than the Planck scale. In this case, it is natural to expect that the proton-decay dimension five operators are suppressed only by this low fundamental scale and hence the proton decay rate is considerably enhanced. To survive the present experimental bound, one needs to assign the negative U(1) charges to the Higgs doublets in this case. Arguments on the branching ratios remain unchanged.

Another important assumption we have made is the degeneracy of the squark and slepton masses in different generations. Let us consider the case where the superparticles of the first two generations are heavy, which is realized when an anomalous U(1) flavor symmetry also mediates SUSY breaking. In this scheme, as far as the scalar tops $\tilde{t}_{L,R}$ and the right-handed scalar tau $\tilde{\tau}_R$ are light like in the case of GUTs, the amplitudes of the LLLL operators are greatly suppressed while those of the RRRR operators remain unsuppressed [38, 39]. Thus, one finds that $p \to K^+\bar{\nu}$ tends to dominate over K^0e^+, $K^0\mu^+$ modes. On the contrary, when only $\tilde{t}_{L,R}$ are light, the RRRR operators are suppressed to the same extent of the LLLL operators. Therefore, the discussions about the branching ratios are not altered.

Acknowledgments

This work was supported in part by the Grant-in-aid from the Ministry of Education, Culture, Sports, Science and Technology, Japan, priority area (#707) "Supersymmetry and unified theory of elementary particles," and in part by the Grants-in-aid No.11640246 and No.12047201.

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