R-Parity Violation: 
Origin of $\mu$-Term and Other Consequences

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Abstract

We propose a new mechanism in which the generation of the supersymmetric $\mu$-term as well charged lepton masses is closely tied to R-parity violation involving heavy vector-like families. A realistic example based on $SU(3)_c \times SU(2)_L \times U(1)_Y$, supplemented by the symmetry $R \times U(1)$, is presented, where $R$ ($U(1)$) denotes a continuous $R$ (flavor) symmetry. In addition to the $\mu$-term, the charged fermion mass hierarchies and mixings, as well as baryon number conservation are also nicely explained. Bilinear $R$-parity violating coupling involving the first generation gives rise to neutrino mass relevant for the small angle $\nu_e - \nu_s$ MSW oscillations, where $\nu_s$ denotes a sterile state. The atmospheric neutrino puzzle is resolved via maximal mixing angle $\nu_\mu - \nu_\tau$ oscillations. The decay of the lightest neutralino (LSP) and leptogenesis are briefly discussed.

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1 Introduction

The well known $\mu$ problem in supersymmetric (SUSY) theories is related to the so-called naturalness issue, namely why (how) a given coupling (or mass) is tiny or zero, without any apparent symmetry reasons. It is unclear how the magnitude of $\mu$ is in the 100 GeV (or so) range and not on the order of $M_{\text{Planck}} (M_P)$. If by some discrete and/or continuous symmetries the $\mu$-term is forbidden at the renormalizable level, its origin can be explained either through higher order non-renormalizable operators or non-minimal Kähler potential. The latter case works within a framework in which SUSY breaking occurs through $N = 1$ (supergravity)SUGRA, and the contribution to the $\mu$-term is $\sim m_{3/2} (\equiv \text{gravitino mass} \sim 10^3 \text{ GeV})$, while other potentially large contributions are absent from the theory with the help of suitable symmetries. However, if SUSY breaking arises through gauge mediation, the induced $\mu$-term through this mechanism will be heavily suppressed, since the gravitino mass $\sim m_S \frac{M_M}{M_P}$ ($m_S \sim 10^3 \text{ GeV}$ is a SUSY scale and $M_M$ is a messenger mass $\lesssim 10^{15} \text{ GeV}$ in order to solve supersymmetric flavor problem). In this case an alternative source for the generation of the $\mu$-term with the desired magnitude is needed. Mechanisms for $\mu$-term generation within the gauge mediated SUSY breaking scenarios were suggested in refs. [6]. A different possibility was recently discussed in [7] where the $\mu$-term arises from a new interaction at the TeV scale.

In this paper we suggest a new mechanism for $\mu$-term generation whose origin is related to ‘matter’ (or R-parity) parity violation, and which leads to some interesting phenomenological consequences. The proposed mechanism turns out to be quite general and can be used to build a variety of realistic models. We begin our considerations by following the naturalness criteria and assume that the $\mu$-term and Yukawa couplings for the charged leptons all vanish at tree level. A crucial role in the generation of both these couplings is played by the vector-like $SU(2)_L$ doublet states $\bar{E} + E$. In the limit when $R$-parity is not violated the $\mu$-term is still zero. By violating $R$-parity in the sector involving the heavy $\bar{E} + E$ superfields and integrating out the latter, the charged lepton masses as well as the $\mu$-term can be generated, which by suitable choice of the parity violating couplings can have the desirable magnitude.

It turns out that the bilinear parity violating operator(s) $h_u l$ can also be generated, so that the effective low energy theory will have some implications different from the minimal supersymmetric standard model (MSSM). For instance, the lightest neutralino (LSP) becomes unstable and another candidate for cold dark matter should be found. Also, one or more of the neutrinos can acquire tree-level mass which should be properly suppressed (see later) in order to be phenomenologically viable.

After presenting the mechanism we indicate the conditions which should be satisfied and outline some clues which can help realize it. We then turn to a specific example and consider a supersymmetric standard model in which SUSY is broken through minimal
$N = 1$ SUGRA and $R$-parity is replaced with the symmetry $\mathcal{R} \times \mathcal{U}(1)$, where $\mathcal{R}$ denotes a continuous abelian $R$-symmetry and $\mathcal{U}(1)$ is an anomalous flavor symmetry. The role of $\mathcal{R} \times \mathcal{U}(1)$ symmetry is three fold. First, it forbids the (direct) $\mu$-term which is generated only through the exchange of heavy states. Second, the $\mathcal{U}(1)$ symmetry allows the possibility of naturally understanding the hierarchies of fermion masses and mixings. Finally, the $\mathcal{R} \times \mathcal{U}(1)$ symmetry also implies baryon number conservation, including higher dimensional operators. We note that in theories with $Z_2$ $R$-parity such as MSSM, the dimension five baryon number violating operators induce unacceptably fast nucleon decay unless some mechanism for their suppression is employed.

In our example $R$-parity (embedded in $\mathcal{R} \times \mathcal{U}(1)$) is violated only in the sector of the first lepton family. The smallness of tree level neutrino mass is also guaranteed by the $\mathcal{R} \times \mathcal{U}(1)$ symmetry. For the generation of lepton masses we introduce three pairs of $\bar{E} + E$ states, while the down quark Yukawa couplings emerge through the exchange of three pairs of $\bar{D}^c + D^c$. It is worth noting that the $\bar{E}$, $D^c$ and $\bar{E}$ states constitute complete $\bar{5} + 5$ multiplets of the $SU(5)$, and because of this the MSSM unification of the three gauge couplings is retained in our model. The resolution of the atmospheric and solar neutrino puzzles requires the introduction of a sterile neutrino state, which is kept light by exploiting the $\mathcal{R} \times \mathcal{U}(1)$ symmetry [9, 10]. The atmospheric neutrino deficit is due to maximal $\nu_\mu - \nu_\tau$ mixing, while the solar neutrino anomaly is resolved through the small angle $\nu_e - \nu_s$ MSW oscillations.

Nearly degenerate right handed neutrino states, which we invoke in the neutrino sector, create a lepton asymmetry through their decays, with the CP asymmetry resonantly enhanced due to the mass degeneracy. This can explain the baryon asymmetry (which will be created from lepton asymmetry during the electroweak phase transition) of our Universe.

2 Mechanism for $\mu$-Term Generation

Let us consider the lepton sector and assume that the charged lepton masses are generated through the exchange of some heavy states. For demonstration of the mechanism we will first consider the case of one generation. The extension to all three generations will be straightforward. We supplement the standard $e^c$, $l$ states with vector-like $\bar{E} + E$ pair, where $E$ has the same transformation properties as $l$, and $\bar{E}$ is conjugate to $E$. Suppose that the direct coupling $e^c l h_d$ vanish by symmetry reasons, and consider the superpotential

$$W = \lambda e^c E h_d + M \bar{E} l + M_E \bar{E} E,$$  \hspace{1cm} (1)

where $\lambda$ is a dimensionless coupling, and $M$, $M_E$ are (heavy) mass scales. Consider the mass matrix

$$W = \lambda e^c E h_d + M \bar{E} l + M_E \bar{E} E,$$  \hspace{1cm} (1)

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where $\lambda$ is a dimensionless coupling, and $M$, $M_E$ are (heavy) mass scales. Consider the mass matrix
Assuming $M \lesssim M_E$, the states $\overline{E} + E$ can be integrated out to yield a lepton mass

$$m_e \simeq \lambda \frac{M}{M_E} h_d .$$

(3)

It is clear that (3) is valid below $M_E$.

The couplings in (1), (2) respect R-parity, if $\overline{E} + E$ are treated as ‘matter’ superfields. For a moment let us assume that we do not have a priori R-parity in the theory and include the following bilinear terms

$$W' = \tilde{m} h_d + m E h_u ,$$

(4)

where $m, \tilde{m}$ are some mass scales which, by assumption, obey the relations:

$$m, \tilde{m} \lesssim M_E , \quad M \lesssim \tilde{m} .$$

(5)

As we will see more precisely below, the first two relations in (5) imply that the physical ‘light’ electroweak Higgs doublets mainly reside in $h_d, h_u$, while the last relation of (5) indicates that the ‘light’ physical left-handed lepton doublet state resides mainly in $l$. If the relations in (5) do not hold, one can redefine the appropriate superfields, so that the conditions in (5) can be taken without loss of generality.

Taking into account the couplings in (1), (4), consider the mass matrix which would be relevant for the generation of both the lepton mass and the $\mu$-term as well (we assume that the direct $\mu$-term is also forbidden by symmetry reasons). Since the states $h_d$ and $E$ mix, the latter will develop a vacuum expectation value (VEV) which must be taken into account during the analysis. The matrix takes the form

$$
\begin{pmatrix}
  l & h_d & E \\
  0 & \lambda h_d & E \\
  0 & M & \tilde{m} \\
  0 & M & M_E \\
\end{pmatrix} .
$$

(6)

Taking account of (5), the states $\overline{E} + E$ can be integrated out which generates the $\mu$-term

$$\mu \simeq \frac{m \tilde{m}}{M_E} .$$

(7)
The scales $m, \tilde{m}, M_E$ should be chosen to obtain a suitable value for $\mu$ (we will see later how symmetries can help us achieve this). The lepton mass will be generated after electroweak symmetry breaking and equals

$$m_e \simeq \frac{\det M}{\mu M_E} = \frac{\lambda M}{\tilde{m}} M_E.$$  

(8)

From (8) one can easily verify that

$$E \supset \frac{\tilde{m}}{M_E} h_d.$$  

(9)

Taking into account (9), from (8) we recover (3).

To summarize, by introducing $\bar{E} + E$ states in the theory and including appropriate $R$-parity violating couplings, together with lepton masses one also can generate the $\mu$-term. The mechanism suggested above is quite general and can be used to construct a variety of models. Before constructing a realistic example, let us outline the conditions which must be satisfied to avoid conflict with phenomenology.

Note that together with the $\mu$-term, integration of $\bar{E} + E$ states also leads to the bilinear $R$-parity violating term

$$\mu_l h_u,$$  

(10)

where, according to (6),

$$\mu_e \simeq \frac{M m}{M_E}.$$  

(11)

This term, in general, can create a non-zero sneutrino VEV, which will lead to a neutrino mass through mixings with neutralinos [11]-[14]. Without any additional mechanism of suppression, the neutrino mass is expected to be in the 100 GeV range!

From (7), (11) the generalized supersymmetric $\mu$-terms’ have the form

$$W_\mu = (\mu_e l + \mu h_d) h_u,$$  

(12)

and one may think that, after suitable rotation of the doublets, only one combination $h'_d$ will have coupling with $h_u$. However, this is not so, because after SUSY breaking there also emerge the well known soft terms $(m^2_s)_{ij}$ and $A, B \ (i, j$ refer to superfields with identical transformation properties under $SU(3)_c \times SU(2)_L \times U(1)_Y$). These soft terms, in general, are not universal and proportional (to $W$) and therefore, in the minimization of the Higgs potential the scalars of both $l$ and $h_d$ superfields participate. As was shown in [11, 13], the neutrino mass will vanish if
\[(m^2_{\nu})_{ij} = m^2_\delta \delta_{ij}, \quad B_{\mu_i} \sim \mu_i. \] (13)

The neutrino mass is expected to be \([11]-[14]\)
\[m_\nu \sim O(100 \text{ GeV}) \sin^2 \xi, \] (14)

where the misalignment parameter
\[\sin \xi = \frac{B_{\mu_e} \langle h_d \rangle - B_\mu \langle l \rangle}{(\langle h_d \rangle^2 + \langle l \rangle^2)^{1/2}(B_{\mu_e}^2 + B_\mu^2)^{1/2}}. \] (15)

In the case of (13) one has the alignment \([11],[13]\)
\[\frac{\mu_e}{\mu} = \frac{B_{\mu_e}}{B_\mu} = \frac{\langle l \rangle}{\langle h_d \rangle}, \] (16)

so that \(\sin \xi = 0\) and the neutrino mass vanishes.

An alternative way for suppressing the neutrino mass is to have a hierarchy between \(\mu_e\) and \(\mu\), \(B_{\mu_e}\) and \(B_\mu\) \([11]-[13],[15]\)
\[\mu_e \ll \mu, \quad B_{\mu_e} \ll B_\mu, \quad \mu_e \mu^{-1} \sim \frac{B_{\mu_e}}{B_\mu}. \] (17)

In this case \(\sin \xi \sim \frac{\mu_e}{\mu}\), and the suppression factor for the neutrino mass will be \((\frac{\mu_e}{\mu})^2\).

The alignment in (13) can be achieved when SUSY is broken through minimal \(N=1\) SUGRA so that (16) is realized. For this case, with the mechanism discussed above, the ‘\(B_\mu\)-terms’ are generated after inclusion of the soft terms
\[V_{SB} = A_\lambda e^c E h_d + B_m h_u E + B_M E l + B_{\tilde{m}} E h_d + B_{M_E} E E, \] (18)

where the scalar components of appropriate superfields are assumed in (18), and at \(M_P,\)
\[A_\lambda = m_{3/2} A \lambda, \quad B_m = m_{3/2} m, \quad B_M = m_{3/2} M, \]
\[B_{\tilde{m}} = m_{3/2} \tilde{m}, \quad B_{M_E} = m_{3/2} M_E. \] (19)

Taking into account (1), (4), (18) it is easy to verify that after integration of \(E + \bar{E}\) states, the \(B_\mu\)-terms are generated
\[B_\mu = \frac{m B_{\tilde{m}} + \tilde{m} B_m}{M_E}, \quad B_{\mu_e} = \frac{m B_M + M B_{\tilde{m}}}{M_E}, \] (20)

and, using (19), we see that \(B_\mu\)-terms are aligned with the \(\mu\)-terms as in (14). However, both universality and proportionality (19) only hold at \(M_P\). The alignment in MSSM is violated due to renormalization effects and one expects \([14]\)
\[
\sin \xi \sim \frac{Y_b^2}{16\pi^2} \ln \frac{M_Z}{M_P},
\]
where MSSM in parenthesis is inserted to remind the reader that the estimate only holds if the field content is identical with that of MSSM.

In our model there is a new source for misalignment. Since \(\mu\) and \(B_\mu\)-terms are generated at \(M_E\), the states \(\overline{E} + E\) (and also the states \(\overline{D}^c + D^c\), which would generate down quark masses as in sect. (3.1) below) will provide additional contribution to the misalignment through renormalization between \(M_P\) and \(M_E\). Let us estimate this effect. The renormalization group (RG) equations for the relevant parameters read

\[
16\pi^2 \frac{d m_i}{dt} = m_i \left( \lambda_{ji} \lambda_{ji} + 3 \text{tr}(Y_u Y_u^T) - 3g^2 - g'^2 \right)
\]

\[
16\pi^2 \frac{d \bar{m}_i}{dt} = \bar{m}_i \left( \text{tr}(\lambda \lambda^T) + 3 \text{tr}(\lambda_D \lambda_D^T) - 3g^2 - g'^2 \right)
\]

\[
16\pi^2 \frac{d M_i}{dt} = M_i \left( -3g^2 - g'^2 \right)
\]

\[
16\pi^2 \frac{d M_E}{dt} = M_E \left( \lambda_{ji} \lambda_{ji} - 3g^2 - g'^2 \right).
\]

where \(\lambda_D\) denotes couplings which will appear if the down quark masses are also induced by integration of colored states circulating in the loops. Taking into account (17) (11), from (22) we obtain

\[
16\pi^2 \frac{d}{dt} \ln \frac{\mu_e}{\mu} = -\text{tr}(\lambda \lambda^T) - 3 \text{tr}(\lambda_D \lambda_D^T).
\]

Assuming that \(\lambda, \lambda_D \ll 1\), from (23) one obtains

\[
\left. \frac{\mu_e}{\mu} \right|_{M_P} - \left. \frac{\mu_e}{\mu} \right|_{M_E} \approx \frac{1}{16\pi^2} [\text{tr}(\lambda \lambda^T) + 3 \text{tr}(\lambda_D \lambda_D^T)] \ln \frac{M_E}{M_P}.
\]

Clearly, analogous relations also hold for \(B_\mu\)-terms because the \(A_\lambda\) couplings are suppressed like \(\lambda\) (see (19)). Therefore, the expected misalignment is estimated to be

\[
\sin \xi \sim \frac{1}{16\pi^2} [\text{tr}(\lambda \lambda^T) + 3 \text{tr}(\lambda_D \lambda_D^T)] \ln \frac{M_E}{M_P}.
\]

Note that in (25), \(\lambda, \lambda_D\) appear because they are the only Yukawa couplings between \(M_E\) and \(M_P\) which can induce misalignment. If they are small enough, the neutrino mass will still be suppressed.
The second mechanism for suppressing the neutrino mass requires a hierarchical structure \[^{[17]}\], and can be realized through the flavor symmetries \[^{[11]-[13], [15]}\]. It is possible, of course, that these two suppression mechanisms of tree level induced neutrino mass work together, in which case

\[
\sin \xi \sim \frac{\mu_e}{\mu} \frac{1}{16\pi^2} \left[ \text{tr}(\lambda\lambda^T) + 3\text{tr}(\lambda_D\lambda_D^T) \right] \ln \frac{M_E}{M_P} .
\]

(26)

In order to have neutrino mass in the range \(< 0.1 \text{ eV}\), one should have \(\sin \xi \lesssim 3 \cdot 10^{-7}\) (see \[^{[14]}\]).

Indeed, these two mechanisms are simultaneously present in the model presented below.

3 The Model

Consider the supersymmetric standard model with \(\mathcal{R} \times \mathcal{U}(1)\) symmetry and no \(R\)-parity a priori. Under \(\mathcal{R}\), \(W \to e^{iR}W\), \(\phi_i \to e^{iR_i}\phi_i\), where \(R_i\) is the \(R\)-charge of the superfield \(\phi_i\). \(\mathcal{U}(1)\) is a flavor symmetry which is anomalous. As emphasized earlier, \(\mathcal{R} \times \mathcal{U}(1)\) will be crucial for the realization of the mechanism presented above, and for a natural explanation of the hierarchies between fermion masses and mixings. The \(\mathcal{R}\) and \(\mathcal{U}(1)\) symmetry breaking scales also play a crucial role in our considerations.

Let us start with the description of \(\mathcal{R} \times \mathcal{U}(1)\) symmetry breaking. We introduce the singlet superfields \(Z, \bar{Z}, X\) with the following transformation properties under \(\mathcal{R}\) and \(\mathcal{U}(1)\):

\[
\mathcal{R} : \quad R_W = R , \quad R_Z = \frac{2R}{5} , \quad R_{\bar{Z}} = -\frac{R}{5} , \quad R_X = 0 ,
\]

\[
\mathcal{U}(1) : \quad Q_Z = q , \quad Q_{\bar{Z}} = -q , \quad Q_X \neq 0 .
\]

(27)

The charges \(q, Q_X\) are not fixed for the time being. However, let us note that the neutrino sector helps fix these charges (see \[^{(64)}\]) in such a way that the single allowed term in the scalar superpotential involving the \(Z, \bar{Z}, X\) superfields, is

\[
W_s = M_P^3 \left( \frac{Z\bar{Z}}{M_P^2} \right)^5 .
\]

(28)

In the unbroken SUSY limit the VEVs \(\langle Z \rangle, \langle \bar{Z} \rangle\) are zero. After SUSY breaking through minimal \(N = 1\) SUGRA, together with the soft terms

\[
V_m = m^2_{3/2} (|Z|^2 + |\bar{Z}|^2) ,
\]

(29)
one finds:
\[
\frac{|\langle Z \rangle|}{M_P} = \frac{|\langle \overline{Z} \rangle|}{M_P} \equiv \epsilon_G \sim \left( \frac{m_{3/2}}{M_P} \right)^{1/8} \approx 10^{-2},
\]
(30)
where we have taken \( m_{3/2} = 10^3 \) GeV, \( M_P = 2.4 \cdot 10^{18} \) GeV. The scale of \( \mathcal{R} \) symmetry breaking is therefore close to the GUT scale (\( \sim 10^{16} \) GeV). Since the \( \mathcal{U}(1) \) symmetry is anomalous, the Fayet-Iliopoulos term
\[
\xi \int d^4\theta V_A
\]
will be generated \([16]\), where, in string theory \([17]\)
\[
\xi = \frac{g^2_A M_P^2}{192\pi^2} \text{Tr} Q.
\]
(32)
The \( D_A \)-term will have the form
\[
\frac{g_A^2}{8} D_A^2 = \frac{g_A^2}{8} \left( \Sigma Q_i |\phi_i|^2 + \xi \right)^2,
\]
(33)
where \( Q_i \) is the ‘anomalous’ charge of \( \phi_i \) superfield. With opposite signs of \( \xi \) and \( Q_X \), the cancellation of (33) fixes a non-zero VEV for the scalar component of \( X \),
\[
\langle X \rangle = \left( -\frac{\xi}{Q_X} \right)^{1/2}.
\]
(34)
We will assume that the scale of \( \mathcal{U}(1) \) symmetry breaking is
\[
\frac{\langle X \rangle}{M_P} \equiv \epsilon \approx 0.22.
\]
(35)
In refs. \([18]\) the anomalous \( \mathcal{U}(1) \) symmetry was considered as a mediator of SUSY breaking, while in refs. \([19]\) the anomalous Abelian symmetries were exploited as flavor symmetries for a natural understanding the hierarchies of fermion masses and mixings. The parameter \( \epsilon \) is an important expansion parameter in our scheme. Below we will express the magnitudes of Yukawa couplings and CKM matrix elements in terms of \( \epsilon \) and \( \epsilon_G \) (see (30)).

### 3.1 \( \mu \)-Term, Charged Fermion Masses and Mixings and Related Issues

We start our considerations with the lepton sector and introduce an additional three families of vector-like supermultiplets \( \overline{E} + E \). These states, together with \( e^c, l, h_d, h_u \), have flavor-universal \( \mathcal{R} \) charges:
\[ R_{ec} = \frac{6R}{5}, \quad R_l = -\frac{9R}{10}, \quad R_{\bar{E}} = \frac{7R}{10}, \quad R_E = \frac{3R}{10}, \]

\[ R_{h_d} = -\frac{R}{2}, \quad R_{h_u} = \frac{17R}{10}, \]

while the \( U(1) \) assignment has flavor dependent structure:

\[ Q_{e_1} = 2q + 4Q_X, \quad Q_{e_2} = Q_{e_3} = 2q + \frac{11}{2}Q_X, \quad Q_{l_1} = -3q - 3Q_X, \]

\[ Q_{l_2} = -3q - \frac{3}{2}Q_X, \quad Q_{l_3} = -3q + \frac{1}{2}Q_X, \quad Q_{E_1} = -6Q_X, \]

\[ Q_{E_2} = Q_{E_3} = -\frac{15}{2}Q_X, \quad Q_{\bar{E}_1} = -2Q_X, \quad Q_{\bar{E}_2} = \frac{3}{2}Q_X, \]

\[ Q_{\bar{E}_3} = -\frac{1}{2}Q_X, \quad Q_{h_d} = -2q, \quad Q_{h_u} = 6q + 6Q_X. \]

With the prescriptions (36), (37), and taking into account (64), one observes that the direct coupling \( h_d h_u \) is forbidden to all orders, and the tree level Yukawa couplings \( e^c l h_d \) also vanish. The presence of \( \bar{E} + E \) states is therefore crucial.

As we see from (37), the states of the second and third generations have non-integer \( Q_X \) charges and therefore will not participate in the type of couplings in (4), while the states from the first family will be relevant for the generation of the \( \mu \)-term. The relevant couplings have the following matrix representation:

\[
\begin{pmatrix}
  l_1 & h_d & E_1 \\
  0 & (\frac{X}{M_P})^2 E & (\frac{X}{M_P})^2 h_d \\
  0 & 0 & Z (\frac{Z}{M_P})^7 \\
\end{pmatrix}
\]

(38)

Substituting in (38) the VEVs of appropriate superfields (30), (35) and comparing (38) with (3), from the expressions (3), (10), (11) we obtain

\[ h_u(\mu_0 h_d + \mu_1 l_1), \quad \mu \approx M_P \frac{\epsilon G^{10}}{\epsilon_6} \approx 200 \text{ GeV}, \]

\[ \mu_1 \sim \epsilon G^3 \mu, \]

(39)

\[ \lambda_e \sim \frac{\epsilon G^3}{\epsilon}, \]

(40)
We see that the $\mu$-term has just the desired magnitude. Furthermore, from (40) one can verify that the MSSM parameter $\tan\beta$ is close to unity. The bilinear $R$-parity violating $\mu_1$-term in (39) will cause the LSP to be unstable [20], so that an alternative candidate for cold dark matter must be found.

Next we consider the couplings which are relevant for the two heavier generations:

$$
\begin{pmatrix}
 l_2 & l_3 & E_2 & E_3 \\
 0 & 0 & \left(\frac{X}{M_P}\right)^2 h_d & \left(\frac{X}{M_P}\right)^2 h_d \\
 0 & 0 & \left(\frac{X}{M_P}\right)^2 h_d & \left(\frac{X}{M_P}\right)^2 h_d \\
 M_P \left(\frac{Z}{M_P}\right)^3 & 0 & M_P \left(\frac{X}{M_P}\right)^6 & M_P \left(\frac{X}{M_P}\right)^6 \\
 \frac{X^2}{M_P} \left(\frac{Z}{M_P}\right)^3 & M_P \left(\frac{Z}{M_P}\right)^3 & M_P \left(\frac{X}{M_P}\right)^8 & M_P \left(\frac{X}{M_P}\right)^8 \\
\end{pmatrix}
$$

(41)

Integration of the heavy $(E + E)_{2,3}$ states yields

$$
e_2^c \begin{pmatrix} l_2 & l_3 \\ \epsilon^2 & 1 \end{pmatrix} e_3^c h_d.
$$

(42)

From (42), (40) we find

$$
\lambda_\tau \sim \frac{\epsilon_2^3}{\epsilon_6} \sim 10^{-2}, \quad \tan \beta \sim 1,
$$

$$
\lambda_e : \lambda_\mu : \lambda_\tau \sim \epsilon^5 : \epsilon^2 : 1,
$$

(43)

which is indeed the desirable hierarchical structure for the Yukawa couplings of the charged leptons.

Turning to the quark sector, for generating the down quark masses we introduce three pairs of $D^c + D^c$. With the transformation properties:

$$
R_q = R_{e^c}, \quad R_{d^c} = R_t, \quad R_{D^c} = R_E, \quad R_{\bar{T}^c} = R_{\bar{E}}
$$

(44)

$$
Q_{q_1} = 2q + \frac{5}{2} Q_X, \quad Q_{q_2} = 2q + \frac{7}{2} Q_X, \quad Q_{q_3} = 2q + \frac{11}{2} Q_X, \quad
Q_{d_1^c} = -3q - \frac{1}{2} Q_X, \quad Q_{d_2^c} = Q_{d_3^c} = -3q + \frac{1}{2} Q_X, \quad
Q_{D_1^c} = -\frac{15}{2} Q_X, \quad Q_{\bar{D}_1^c} = -\frac{1}{2} Q_X,
$$

(45)

the mass matrix relevant for down quark masses has the form
\[
\frac{q}{D^c} \left( \begin{array}{cc}
d^c & D^c \\
0 & \hat{A} h_d \\
\hat{M}_{D^c} & \hat{M}_{D^c}
\end{array} \right),
\]

where
\[
\hat{A} = \left( \begin{array}{ccc}
e^3 & e^3 & e^3 \\
e^2 & e^2 & e^2 \\
1 & 1 & 1
\end{array} \right) e^2, \quad \hat{M}_{D^c} = \left( \begin{array}{ccc}
e & 1 & 1 \\
e & 1 & 1 \\
e & 1 & 1
\end{array} \right) \frac{Z^3}{M_P},
\]

\[
\hat{M}_{D^c}^H = M_P \left( \frac{X}{M_P} \right)^8 \alpha^{ij}
\]

(\(\alpha^{ij}\) are dimensionless couplings of order unity). Integrating out the heavy decoupled states gives
\[
\hat{m}_d = A \hat{M}_{D^c}^{-1} \hat{M}_{D^c} h_d \simeq q_1 \left( \begin{array}{ccc}
d_1^c & d_2^c & d_3^c \\
e^4 & e^3 & e^3 \\
e & 1 & 1
\end{array} \right) \frac{e^2 h_d}{e^6},
\]

which upon diagonalization yields
\[
\lambda_b \sim \frac{e^3 G}{e^6} \sim 10^{-2}, \quad \lambda_d : \lambda_s : \lambda_b \sim e^4 : e^2 : 1,
\]

the desired hierarchies between the three families of down quarks.

Before discussing the up quark sector, let us note that the heavy decoupled states (\(E + E\) doublets and \(D^c + D^c\) triplets) have masses
\[
m_{d_1} \simeq m_{d_3} \simeq M_P e^8, \quad m_{d_2} \simeq M_P e^6,
\]

\[
m_{t_1} \simeq m_{t_2} \simeq m_{t_3} \simeq M_P e^8,
\]

and constitute three \(\bar{5} + 5\) states of \(SU(5)\). This allows the possibility to retain the successful unification of the three gauge couplings, and also obtain an improved value for \(\alpha_s(M_Z)\)
\[
\alpha_s^{-1} = (\alpha_s^0)^{-1} - \frac{9}{14\pi} \ln \frac{m_{t_1} m_{t_2} m_{t_3}}{m_{d_1} m_{d_2} m_{d_3}},
\]

where \(\alpha_s^0\) is the strong coupling, calculated at \(M_Z\) in minimal SUSY \(SU(5)\). Taking \((\alpha_s^0)^{-1} = 1/0.126 [21]\) and using (51), from (52) we obtain \(\alpha_s = 0.117\), which is in excellent agreement with world average value [22].
Turning to the up quark sector, we prescribe the following transformation properties to the \( u^c \) states

\[
R_{u^c} = -\frac{19}{10} R, \quad Q_{u^c} = -8q - \frac{29}{2} Q_X ,
\]

\[
Q_{u^c_1} = -8q - \frac{25}{2} Q_X , \quad Q_{u^c_2} = -8q - \frac{23}{2} Q_X . \tag{53}
\]

The up quark mass matrix will have the structure:

\[
\tilde{m}_u \approx \begin{pmatrix}
q_1 & u^c_1 & u^c_2 & u^c_3 \\
q_2 & \epsilon^6 & \epsilon^4 & \epsilon^3 \\
q_3 & \epsilon^3 & \epsilon & 1
\end{pmatrix}
\]

whose diagonalization yields

\[
\lambda_t \sim 1, \quad \lambda_u : \lambda_c : \lambda_t \sim \epsilon^6 : \epsilon^3 : 1 . \tag{55}
\]

From (49) and (54) one can also estimate the CKM matrix elements

\[
V_{us} \sim \epsilon , \quad V_{cb} \sim \epsilon^2 , \quad V_{ub} \sim \epsilon^3 . \tag{56}
\]

which is in very good agreement with observations.

It is worth noting that since the states \( d^c, u^c \) have non-integer \( Q_X \) charges, the baryon number violating trilinear couplings \( u^c d^c d^c \) are forbidden. This is a consequence of the fact that in quark sector, due to \( U(1) \) symmetry, \( R \)-parity emerges automatically.

The Planck scale suppressed baryon number violating \( d = 5 \) operators

\[
\frac{1}{M_P} \bar{q}q l , \quad \frac{1}{M_P} \bar{q}q E , \quad \frac{1}{M_P} \bar{q}q h_d , \\
\frac{1}{M_P} u^c u^c d^c e^c , \quad \frac{1}{M_P} u^c d^c e^c , \tag{57}
\]

are also forbidden by the \( R \) symmetry \([10]\). Thanks to the \( R \times U(1) \) symmetry, baryon number is conserved in our scheme \([23]\).

As far as lepton number violation is concerned, the couplings \( e^c l l , e^c E E \) and \( qD^c E \) are also forbidden by \( R \times U(1) \). The \( qd^c l \) type operator has at least the suppression factor \( (\frac{Z}{M_P})^4 \), which satisfies all phenomenological bounds \([25]\) and will not have any significant contributions to neutrino masses (that emerge radiatively through this coupling).

\[\text{4 Different mechanism for baryon number conservation in } R \text{-parity violating GUT scenario was considered in } [24], \text{ where a missing VEV solution of the GUT adjoint was applied.}\]
Some other allowed operators are

\[ \frac{Z}{M_P} \Gamma_{ijk} e_i^c E_j l_k , \quad \frac{Z}{M_P} \Gamma'_{ijk} q_i d_j^c E_k , \quad \frac{Z}{M_P} \Gamma''_{ijk} q_i D_j^c l_k \]  

(58)

(\Gamma, \Gamma' and \Gamma'' are family dependent couplings), and they all involve the decoupled heavy states. They give rise to the lepton number rotating operators

\[ \lambda_{ijk} e_i^c l_j^c l_k , \quad \lambda'_{ijk} q_i d_j^c l_k . \]  

(59)

Taking into account (38), (41), (46)-(48) one can verify that

\[ E_3 \supset \frac{e^3}{\epsilon^8} \left( e^2 l_2 + l_3 \right) , \quad E_2 \supset \frac{e^3 G}{\epsilon^6} l_2 , \quad E_1 \supset \frac{e^3 G}{\epsilon^3} l_1 , \]

\[ D_i^c \supset \frac{e^3}{\epsilon^8} \left( \epsilon d_1^c + d_{2,3}^c \right) . \]  

(60)

The \( \Gamma \) factors in (58) are expressed through appropriate powers of \( X \), depending on the appropriate superfield and can be selected from the prescriptions (37), (45). From this, taking into account (58), (60), the non-zero \( \lambda \) and \( \lambda' \) suppression factors are (let us note that \( \lambda_{ijk} = -\lambda_{ikj} \))

\[ \lambda_{123} \sim \lambda_{213} \sim \lambda_{312} \sim \lambda'_{321} \sim \lambda'_{331} \sim \frac{e^4 G}{\epsilon^3} , \]

\[ \lambda_{212} \sim \lambda_{312} \sim \lambda'_{221} \sim \lambda'_{231} \sim \frac{e^4 G}{\epsilon^3} , \]

\[ \lambda'_{111} \sim \epsilon^4 G , \quad \lambda'_{121} \sim \lambda'_{131} \sim \lambda'_{211} \sim \epsilon^4 , \quad \lambda'_{311} \sim \frac{e^4 G}{\epsilon^2} . \]  

(61)

All other trilinear lepton number violating couplings which are not presented in (61) are zero due to \( R \times U(1) \) symmetry. We note that \( \lambda, \lambda' \) are suppressed at the required level, so that the phenomenological bounds [23] obtained from different processes are satisfied. The contributions (through radiative corrections) from these couplings to the neutrino masses are also negligible.

The \( R \)-parity violating bilinear and trilinear couplings make the LSP unstable. In our model the dominant contribution to LSP decay comes from the bilinear \( \mu_1 \)-term in (39). The lifetime for decay into three fermions is given by

\[ \tau_{\chi}^{-1} = \mu_1^2 Z_{\chi h}^2 \left( \frac{1}{4} + \sin^2 \theta_W + \frac{4}{3} \sin^4 \theta_W \right) \frac{G_F m_{\chi}^3}{192 \pi^3} . \]  

(62)

and for \( m_\chi \sim 100 \) GeV, taking into account (39), we obtain \( \tau_{\chi} \sim 10^{-19} \) sec. Therefore, the LSP is cosmologically irrelevant.
3.2 Neutrino Oscillations and Leptogenesis

In this section we investigate the neutrino sector of our model and attempt to accommodate the recent atmospheric [26] and solar [27] neutrino data. Starting with atmospheric neutrinos, let us note that the prescription (37) of $U(1)$ charges for $l_2, l_3$ permit us to realize maximal mixing between $\nu_\mu$ and $\nu_\tau$ through the mechanism described in [28]. Introducing two right handed $N_{2,3}$ states with transformation properties

$$R_{N_{2,3}} = R \frac{5}{9}, \quad Q_{N_2} = -3q - \frac{19}{2} Q_X, \quad Q_{N_3} = -3q - \frac{9}{2} Q_X,$$

with the condition

$$Q_X = -\frac{5}{14} q,$$

the relevant couplings have the desirable textures:

$$\begin{pmatrix} N_2 & N_3 \\ l_2 & e^5 \\ l_3 & e^3 \end{pmatrix} h_u,$$

$$\begin{pmatrix} N_2 & N_3 \\ N_2 & e^5 \\ N_3 & 1 \end{pmatrix} \frac{Z^2}{M_P}.$$  \hspace{1cm} (65)

After integrating out the heavy $N_{2,3}$ states, the neutrino mass matrix for $\nu_\mu, \nu_\tau$ will have the quasi-degenerate form

$$\begin{pmatrix} \nu_\mu & \nu_\tau \\ e^2 & 1 \\ 1 & 0 \end{pmatrix} m,$$

$$m = \frac{e^3 h_u}{M_P \epsilon_G},$$

with

$$m_{\nu_2} \simeq m_{\nu_3} \simeq m \simeq 0.13 \text{ eV}.$$  \hspace{1cm} (67)

For the atmospheric neutrino oscillation parameters we find

$$\Delta m_{23}^2 = 2 m^2 \epsilon^2 \simeq 2 \cdot 10^{-3} \text{ eV}^2,$$

$$\sin^2 2\theta_{\mu\tau} = 1 - \mathcal{O}(\epsilon^4).$$

The resolution of the solar neutrino puzzle in our scenario requires the introduction of a light sterile neutrino state. We have arranged the lepton sector in such a way as to avoid mixing of the first generation with the second and third generations. This was necessary because $R$-parity violation (through bilinear terms) in the sector of heavy generations could create unacceptably heavy neutrinos. We violated $R$-parity in the sector of first generation and generated the terms $\mu, \mu_1$ in [39]. Due to $\mathcal{R} \times U(1)$ there arises the following hierarchy between $\mu$ and $\mu_1$, and therefore between $B_\mu$ and $B_{\mu_1}$ as well.
\[
\frac{\mu_1}{\mu} = \frac{B_{\mu_1}}{B_\mu} \sim \epsilon G e^3 .
\] (69)

The alignment holds at \( M_P \) since we are working within the framework of minimal \( N = 1 \) SUGRA. This alignment is violated due to renormalization between \( M_P \) and mass scales of \( 'E, D' \) states and, taking into account \((25), (38), (41), (47)\), we expect

\[
\frac{\Delta \mu}{\mu} \sim \frac{6 \epsilon^2}{16 \pi^2} \ln \epsilon^8 .
\] (70)

The misalignment parameter is estimated to be

\[
\sin \xi \sim \frac{\mu_1}{\mu} \frac{6 \epsilon^4}{16 \pi^2} \ln \epsilon^8 \sim \epsilon G e^3 \frac{6 \epsilon^4}{16 \pi^2} \ln \epsilon^8 \sim 10^{-7} ,
\] (71)

and the mass of the \( '\nu_e' \) state (see \((14)\)) is given by

\[
m_{\nu_e} \simeq 10^{-3} \text{ eV} .
\] (72)

As we mentioned above, the state \( \nu_e \) does not mix with \( \nu_{\mu, \tau} \) (which in any case are too heavy for the solar neutrino puzzle). We introduce a sterile state \( \nu_s \) with the transformation properties:

\[
R_{\nu_s} = \frac{R}{5} , \quad Q_{\nu_s} = -3q - 26Q_X .
\] (73)

Taking into account \((64)\), the relevant couplings are

\[
W_{\nu_s} = \left( \frac{X}{M_P} \right)^{23} l_1 \nu_s h_u + \frac{Z_2 Z}{M_P^2} \left( \frac{X}{M_P} \right)^{38} \nu_s^2 ,
\] (74)

from which we have

\[
m_{\nu_s} = M_P \epsilon G e^{38} \simeq 2.5 \times 10^{-4} \text{ eV} , \quad m_{\nu_s \nu_e} = \epsilon^{23} h_u \simeq 10^{-4} \text{ eV} .
\] (75)

Note that \( \nu_s \) is kept light (in \((74), (73)\)) by the symmetry \( R \times U(1) \) \( [4, 10] \).

Taking into account \((72), (75)\), for the solar neutrino oscillation parameters we will have

\[
\Delta m_{\nu_e,\nu_s}^2 \simeq m_{\nu_e}^2 \simeq 10^{-6} \text{ eV}^2 ,
\]

\[
\sin^2 2\theta_{es} \simeq 4 \left( \frac{m_{\nu_s \nu_e}}{m_{\nu_e}} \right)^2 \simeq 10^{-2} ,
\] (76)
which explains the solar $\nu_e$ deficit through the small angle MSW oscillations. The tree level induced $\nu_e$ mass (72), whose origin lies in the $R$-parity violating bilinear (39) term, plays a crucial role in (76).

Before concluding, let us note that even though baryon number is perturbatively conserved in our model, the observed baryon asymmetry can be explained by first creating lepton asymmetry [29] through the out of equilibrium decay of the right handed neutrinos $N_{2,3}$. The electroweak sphalerons [30] would partially transform the lepton asymmetry to the observed baryon asymmetry. The out-of-equilibrium condition reads

$$\Gamma \lesssim KH = 1.7K\sqrt{g_*}\frac{T^2}{M_P},$$  

(77)

where $\Gamma$ is the decay rate of $N_{1,2}$ states into leptons, $H$ is Hubble’s constant, $g_*$ is the effective number of massless degree of freedom, and $K = 1 - 10^3$. The heavy neutrino decay rate is

$$\Gamma_N = \frac{(h^+h)_{ii}}{8\pi} M_{N_i},$$  

(78)

where $h$ can be extracted from $hN/N_u$ type couplings of (65). Using (78) and taking into account (53) one can see that (77) is satisfied for $K = 10^3$.

In addition to the out-of-equilibrium condition (77) we need CP violation, which is necessary for generating the asymmetry. In our model the states $N_2$ and $N_3$ are nearly degenerate in mass. This, as shown in refs. [31], can be a very convenient fact because of the resonance enhancement of the CP asymmetry that it creates. On the other hand, the baryon-to-entropy density ratio $Y_B = n_B/s$, which is created from lepton-to-entropy density ratio through sphalerons, equals [32]

$$Y_B \approx -\frac{1}{3} Y_L \approx -\frac{1}{2K} \frac{\delta_{CP}}{g_*}.$$  

(79)

According to (79), for a large CP asymmetry $\delta_{CP}$ in order to have $Y_B \approx 10^{-10}$ one can take larger values for $K$, so that (77) will be more readily satisfied. Therefore, we conclude that all of the conditions [33] for obtaining the baryon asymmetry of the Universe are satisfied within the framework of our scenario.

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