Some phenomenological studies of a $U(1)_R$ — lepton number model with a right handed neutrino
Some phenomenological studies of a $U(1)_R$– lepton number model with a right handed neutrino

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Abstract. We discuss the Higgs sector and neutrino mass and mixing in the context of a supersymmetric model with a $U(1)_R$ symmetry augmented by a single right handed neutrino superfield. The lepton number of the standard model fermions are identified with the negative of their R-charges. The presence of the neutrino Yukawa coupling ($f$) plays a very interesting role in Dirac neutrino mass generation as well as providing a tree level correction to the Higgs boson mass, when $f \sim O(1)$. In order to fit the complete neutrino mass and mixing, a small breaking of the $U(1)_R$ symmetry is considered in the context of anomaly mediated supersymmetry breaking. We also show that when $f \sim 10^{-4}$, the right handed neutrino can be identified with a 7 keV sterile neutrino dark matter as recent studies have speculated.

1. Introduction
Perhaps the most important discovery in high energy physics in recent times is the observation of a neutral boson by the two LHC based experiments CMS [1] and ATLAS [2], widely speculated to be the Higgs boson with mass $\sim 125$ GeV. All the decay channels of this boson is measured with moderate accuracy [3, 4, 5, 6, 7, 8, 9, 10, 11] and the signal strengths corresponding to different Higgs boson decay mode are found to be well in agreement with the standard model (SM) expectations. These results have constrained many beyond standard model (BSM) scenarios. However indirect hints of new physics cannot be ruled out either. In this regard supersymmetry remains one of the most celebrated BSM candidates although the two LHC based experiments have found no excess over the standard model background. These studies have put severe constraints on the superpartner masses. In this light R-symmetric models with Dirac gauginos are well motivated since they can explain the observation of a 125 GeV Higgs boson and provide a suitable dark matter candidate. Neutrino masses and mixing can be explained at the tree as well as at the one loop level. Finally and most importantly the presence of Dirac gluinos can also relax the bound on the first two generations of squark masses.

Motivated by these results we considered a $U(1)_R$ model [12, 13, 14] and extended it with a single right handed neutrino superfield [15]. The R-charges are identified with the negative of their lepton numbers. An important outcome of having R-symmetry is the absence of Majorana gaugino masses. However, Dirac masses for the gauginos are possible and in order to construct such a Dirac mass term, one must include chiral superfields, living in the adjoint representation of the SM gauge group. For example a singlet $\hat{S}$, a triplet $\hat{T}$ under $SU(2)_L$ and an octet $\hat{O}$ under $SU(3)_c$ help to formulate Dirac gaugino masses. The main artifact of this extension with a right
handed neutrino superfield is to have a Dirac neutrino mass at the tree level itself. With some
tuning of the parameters it is quite easy to show that even in the presence of a large neutrino
Yukawa coupling \( f \sim \mathcal{O}(1) \), a suitable small Dirac neutrino mass can be constructed. However
a large \( f \) has other advantages as well. It provides a large tree level contribution \([17]\) to the
Higgs boson mass, thereby requiring negligible radiative corrections.

The R-symmetric case is interesting and should be explored further. Although, in the context
of our work we consider the R-symmetry to be broken mildly. Gravitino mass happens to be
the order parameter of R-breaking in the case of anomaly mediation. The neutrino Yukawa
coupling plays a very interesting role and leads to novel phenomenology. For example large \( f \)
scenario happens to have a bino like neutralino which is the next to lightest neutralino. The
lightest neutralino turns out to be a Majorana neutrino with mass around 0.1 eV. Since the
Higgs boson can decay to a pair of light neutralinos therefore the presence of this light bino like
state enhances the total decay width of the Higgs boson, which constrain the parameter space
severely and put important bounds on certain parameters. In addition, the small \( f \sim \mathcal{O}(10^{-4}) \)
is also important as because it provides a sterile neutrino eigenstate which can be accommodated
as the 7 keV dark matter \([16]\) with appropriate active-sterile mixing.

2. \( U(1)_R \) – lepton number model with a right handed neutrino

Along with the standard MSSM superfields \( \hat{H}_u, \hat{H}_d, \hat{Q}_i, \hat{U}_i^c, \hat{D}_i^c, \hat{L}_i, \hat{E}_i^c \), the model includes a
pair of vector like inert \( SU(2)_L \) doublet superfields. Chiral superfields \( \hat{S}, \hat{T} \) and \( \hat{O} \) are required
to construct ‘supersoft’ Dirac gaugino mass terms \([17]\). In addition we also have the right handed
neutrino superfield \( \hat{N}^c \). The R-charge assignments of the superfields are as follows,

| \( U(1)_R \) | \( \hat{Q}_i \) | \( \hat{U}_i^c \) | \( \hat{D}_i^c \) | \( \hat{L}_i \) | \( \hat{E}_i^c \) | \( H_u \) | \( H_d \) | \( \hat{R}_u \) | \( \hat{R}_d \) | \( \hat{S} \) | \( \hat{T} \) | \( \hat{O} \) | \( \hat{N}^c \) |
|----------------|-------|-------|-------|-------|-------|-------|-------|-------|-------|-------|-------|-------|-------|
| \( \hat{R} \)    | 1     | 1     | 1     | 0     | 2     | 0     | 0     | 2     | 2     | 0     | 0     | 0     | 0     |

**Table 1.** \( U(1)_R \) charge assignments of the chiral superfields.

In a generic R-symmetric model the scalars possess the same R-charge as that of the
superfields, whereas the fermions acquire R-charge one less than that of the superfields. We
identify the lepton number as the negative of their R-charges \([12]\). The R-charge assignments
from table 1 implies the left handed sneutrino vacuum expectation values (vevs) do not violate
lepton number. Therefore, these vevs can be large and are not constrained from the neutrino
Majorana mass. In order to have an invariant action under R-symmetry, the superpotential
must carry R-charge of two units. Without any loss of generality the generic superpotential can
be simplified by choosing a suitable basis in which one of the sneutrinos acquire a vev, in this
case we choose it to be the electron type sneutrino. In addition, both \( \hat{R}_u \) and \( \hat{H}_d \) fields can be
integrated once the coefficient of the bilinear term \( \hat{H}_u \hat{R}_d \) is considered to be very large. The
sneutrino can then play the role of a down type Higgs field. With all these assumptions in place,
the superpotential is of the form

\[
W = y_{ij}^u \hat{H}_u \hat{Q}_i \hat{U}_j^c + \mu_u \hat{H}_u \hat{R}_d + f \hat{L}_a \hat{H}_u \hat{N}^c + \lambda_S \hat{S} \hat{H}_u \hat{R}_d + 2\lambda_T \hat{H}_u \hat{T} \hat{R}_d - M_R \hat{N}^c \hat{S} + W',
\]

and

\[
W' = \sum_{b=2,3} f_{ib}^L \hat{L}_b \hat{L}_b^c + \sum_{k=1,2,3} f_{ik}^L \hat{L}_k \hat{D}_k^c + \sum_{k=1,2,3} \frac{1}{2} \lambda_{23k} \hat{L}_2 \hat{L}_3 \hat{E}_k^c + \sum_{j,k=1,2,3, b=2,3} \lambda_{bjk} \hat{L}_b \hat{Q}_j \hat{D}_k^c,
\]

(2)
where \( a = e \) and \( b = \mu, \tau \). One can always choose a basis in which the scalar component of the superfield \( \tilde{L}_a \) acquires a non zero vev (i.e. \( < \tilde{\nu} > \equiv v_a \neq 0 \)) and the other two sneutrinos do not get any vev. We also note from eq. (2) that such an R-charge assignment enables us to write down a subset of lepton number preserving but trilinear R-parity violating terms in the superpotential.

The soft supersymmetry breaking but \( U(1)_R \) preserving terms in the scalar sector can be generated by a spurious superfield \( X \), defined as \( X = x + \theta^2 F_X \). A nonzero vev of the \( F \) term generates soft supersymmetry breaking terms as written,

\[
V_{soft} = m_H^2 H_u^\dagger H_u + m_{R_d}^2 R_d^\dagger R_d + m_{L_a}^2 \tilde{L}_a \tilde{L}_a + \sum_{b=2,3} m_{L_b}^2 \tilde{L}_b \tilde{L}_b + M_Y^2 \tilde{N}^c \tilde{N}^c + m_{R_i}^2 \tilde{R}_i \tilde{R}_i
\]

\[+ m_3^2 S^\dagger S + 2m_3^2 \text{tr}(T^I T^I) + 2m_3^2 \text{tr}(O^I O) - (b_{\mu L} H_u \tilde{L}_a + \text{h.c.}) + (t_S S + \text{h.c.})
\]

\[+ \frac{1}{2} b_S (S^2 + \text{h.c.}) + b_T (\text{tr}(TT) + \text{h.c.}) + B_O (\text{tr}(OO) + \text{h.c.}). \tag{3}
\]

Finally, the Dirac gaugino masses are generated by a spurious superfield \( W'_a = \lambda_a D' \) where the supersymmetry breaking is of the D-type. The relevant Lagrangian takes the form \([18]\

\[L_{\text{Dirac gaugino}} = \int d^2 \theta \frac{W'_a}{\Lambda} [\sqrt{2} \kappa_1 W_{1a} \hat{S} + 2\sqrt{2} \kappa_2 \text{tr}(W_{2a} \hat{T}) + 2\sqrt{2} \kappa_3 \text{tr}(W_{3a} \hat{O})] + \text{h.c.} \tag{4} \]

3. Scalar sector

In the R-symmetric scenario the CP even scalar squared mass matrix written in the basis \((h_R, \tilde{\nu}_R, S_R, T_R)\) is

\[(M_S^2)_{11} = \frac{(g^2 + g'^2)}{2} v^2 \sin^2 \beta + (f M_{RV} - b \mu L)(\tan \beta)^{-1} + 2\delta \mu v^2 \sin^2 \beta,
\]

\[(M_S^2)_{12} = f v^2 \sin 2\beta + b \mu L - \frac{(g^2 + g'^2 - 2\delta \lambda)}{4} v^2 \sin 2\beta - f M_{RV}, \]

\[(M_S^2)_{13} = 2\lambda_\mu S \nu v \sin \beta + 2\lambda S \nu v T \nu T v \sin \beta \pm \sqrt{2} g M^D_2 v \sin \beta - f M_{RV} \cos \beta,
\]

\[(M_S^2)_{14} = 2\lambda S \nu T v \sin \beta + 2\lambda \mu T \nu v \sin \beta + 2\lambda S \nu T v S \nu S \sin \beta \pm \sqrt{2} g M^D_2 v \sin \beta,
\]

\[(M_S^2)_{22} = \frac{(g^2 + g'^2)}{2} v^2 \cos^2 \beta + (f M_{RV} - b \mu L) \tan \beta + 2\delta \mu v^2 \cos^2 \beta,
\]

\[(M_S^2)_{23} = -\sqrt{2} g M^D_2 v \cos \beta - f M_{RV} \sin \beta,
\]

\[(M_S^2)_{24} = -\sqrt{2} g M^D_2 v \cos \beta,
\]

\[(M_S^2)_{33} = -\mu \lambda S \nu v^2 \sin^2 \beta - \frac{\lambda \mu S \nu v T v^2 \sin^2 \beta}{\nu S} - t_S v + \frac{g' M^D_1 v^2 \cos 2\beta}{\sqrt{2} \nu S} + \frac{f M_{RV} v^2 \sin 2\beta}{2 \nu S},
\]

\[(M_S^2)_{34} = \lambda S \nu T v^2 \sin^2 \beta,
\]

\[(M_S^2)_{44} = -\mu \lambda T \nu v^2 \sin^2 \beta - \lambda S \nu T v S \nu S v^2 \sin \beta - \frac{g M^D_2 v^2}{\sqrt{2} \nu T} \cos 2\beta, \tag{5}
\]

When the singlet and the triplet vevs, \( v_S \) and \( v_T \) are small the corresponding scalar masses becomes very heavy. In such a scenario the scalar mass matrix can be essentially considered to be a \( 2 \times 2 \) mass matrix. The lightest eigenvalue of the CP even scalar mass matrix can be identified with the 125 GeV Higgs boson. One very important outcome of this decoupling limit is the Higgs boson mass gets a contribution at the tree level which is of the following form \([15]\

\[m_H^2 = m_{h^2} + m_{h^2} \cos 2\beta \]
\[
M_h^2 = \left( M_Z^2 \cos^2 2\beta + f^2 v^2 \sin^2 2\beta \right). 
\] (6)

From eq. (6) it is conspicuous that the Higgs boson mass receives substantial contribution at the tree level itself for \( f \sim O(1) \) and for certain regions of \( \tan \beta \). This statement is transparent from figure 1 that \( M_h \sim 125 \) GeV for order one Yukawa coupling and when \( \tan \beta \) is small.

4. Neutrino sector: R preserving case

In the sector of neutral fermions we have a mixing between neutrino, neutral gauginos and neutral higgsinos, i.e. the electron type neutrino \( \nu_e \) mixes with the sterile neutrino \( N^c \). The neutralino mass matrix written in the basis \( \psi^0_+ = (\tilde{b}_0, \tilde{w}_0, \tilde{R}_0, N^c) \) and \( \psi^0_- = (\tilde{S}, \tilde{T}_0, \tilde{H}_u, \nu_e) \)

\[
M_x^D = \begin{pmatrix}
M_1^D & 0 & \frac{g^* v_u}{\sqrt{2}} & -\frac{g^* v_u}{\sqrt{2}} \\
0 & M_2^D & -\frac{g v_u}{\sqrt{2}} & \frac{g v_u}{\sqrt{2}} \\
\lambda_S v_u & \lambda_T v_u & \mu_u + \lambda_S v_S + \lambda_T v_T & 0 \\
M_R & 0 & -f v_u & -f v_u
\end{pmatrix}.
\] (7)

The lightest eigenvalue of this neutralino mass matrix corresponds to the Dirac neutrino mass which we deduce as [15],

\[
m_{\nu_e}^D = \frac{v^3 f g \sin \beta}{\sqrt{2} \gamma M_1^D M_2^D} \lambda_T (M_2^D - M_1^D),
\] (8)

where we assumed the following relations, \( \lambda_T = \lambda_S \tan \theta_W \) and \( M_R = \frac{\sqrt{2} f M_2^D \tan \beta}{g \tan \theta_W} \). We observe from eq. (8) that even in the presence of \( f \sim O(1) \), Dirac neutrino mass can be appropriately small for small \( \lambda_T \sim 10^{-6} \) and nearly degenerate Dirac gaugino masses.
5. Neutrino sector: R breaking case

Recent cosmological observations predict a vanishingly small value for the cosmological constant. To explain this in the framework of a spontaneously broken supergravity theory, the superpotential has to acquire a non-zero vev at the vacuum. The gravitino mass is also proportional to the vev of the superpotential in the vacuum, therefore the gravitino can be thought of as the order parameter of R-breaking. The information of the breaking of R-symmetry would have to be transferred from the hidden sector to the visible sector and we assume anomaly mediation (AMSB) playing the role of this messenger. The breaking of R-symmetry would generate Majorana gaugino masses and trilinear scalar couplings. In the R-breaking scenario the neutralino mass matrix written in the basis \( \tilde{b}, \tilde{S}, \tilde{w}^0, \tilde{T}, \tilde{R}_d^0, \tilde{H}_u^0, N^c, \nu_e \) is given as,

\[
M^M_\chi = \begin{pmatrix}
M_1 & M_1^D & 0 & 0 & 0 & \frac{g' v_u}{\sqrt{2}} & 0 & \frac{g' v_u}{\sqrt{2}} \\
M_1^D & 0 & 0 & 0 & \frac{\lambda_S v_u}{\sqrt{2}} & 0 & M_R & 0 \\
0 & 0 & M_2 & M_2^D & 0 & -\frac{g v_u}{\sqrt{2}} & 0 & \frac{g v_u}{\sqrt{2}} \\
0 & 0 & 0 & M_2^D & 0 & \frac{\lambda_T v_u}{\sqrt{2}} & 0 & 0 \\
0 & \frac{g' v_u}{\sqrt{2}} & 0 & \frac{\lambda_S v_u}{\sqrt{2}} & 0 & \mu_u + \lambda_S v_S + \lambda_T v_T & 0 & -f v_u & 0 \\
0 & M_R & 0 & 0 & 0 & -f v_u & 0 & -f v_u & 0 \\
-\frac{g' v_u}{\sqrt{2}} & 0 & \frac{g v_u}{\sqrt{2}} & 0 & 0 & 0 & 0 & 0 & 0 \\
\end{pmatrix}
\]

The lightest neutralino corresponds to the Majorana neutrino which has a mass as,

\[
(m_\nu)_{\text{Tree}} \simeq -v^2 \left[ \frac{g \lambda_T v^2 (M_2^D - M_1^D) \sin \beta}{M_1 \alpha^2 + M_2 \delta^2} \right]^2,
\]

where

\[
\alpha = \frac{2 M_1^D M_2^D (\mu_u + \lambda_S v_S + \lambda_T v_T) \tan \beta}{g \tan \theta_W} + \sqrt{2} v^2 \lambda_S \tan \beta (M_1^D \sin^2 \beta + M_2^D \cos^2 \beta),
\]

\[
\delta = \sqrt{2} M_1^D v^2 \lambda_T \tan \beta.
\]

The full neutrino mass and mixing can be explained once one loop radiative corrections are taken into consideration as shown in figure 2.\(^1\). The characteristics of the next to lightest neutralino depend on the magnitude of the neutrino Yukawa coupling. Let us now discuss these two cases separately.

5.1. Case with small neutrino Yukawa coupling: \( f \sim \mathcal{O}(10^{-4}) \)

In the limit when \( f \sim \mathcal{O}(10^{-4}) \) the Majorana neutrino mass is given in eq. (10). The sterile neutrino mass computed as \( M_{\tilde{S}}^R \simeq M_1 \frac{f^2 \tan^2 \beta}{g_2^2} \). The active-sterile mixing is estimated as \( \theta_{14}^2 \simeq \frac{(m_\nu)_{\text{Tree}}}{M_{\tilde{S}}^R} \) [16]. With these analytical expressions in hand, we fit a 7 keV sterile neutrino dark matter explaining the emergence of a weak line at 3.5 keV in the X-ray spectra of the Andromeda galaxy and various other galaxy clusters including the Perseus galaxy cluster as observed by XMM-Newton X-ray space observatory [19, 20]. The observed flux and the peak of the energy

\(^1\) For detailed calculations see [14, 15]
Figure 2. One loop diagrams to generate neutrino masses.

can be translated to an active-sterile mixing in the range $2.2 \times 10^{-11} < \sin^2 2\theta_{14} < 2 \times 10^{-10}$. The mass and mixing can be easily conceived in our model. From figure 3 we observe that the Dodelson-Widrow (DW) [21] mechanism of sterile neutrino production is not compatible with such a small active-sterile neutrino mixing as pointed out in refs. [19, 20]. However, resonant production of the sterile neutrinos [22, 23] in the presence of a cosmological lepton asymmetry $L \sim O(10^{-3})$ is capable of producing correct relic density of 0.119 [24].

Further, such a precise measurement of the mass and the mixing of this sterile neutrino provides substantial bound on the model parameters. For example the maximum value of the active-sterile mixing as computed from the recent energy peak at 3.5 keV corresponds to the minimum of $\tan \beta$ and vice-versa. Therefore in figure 4 we show the variation of $\tan \beta$ as a function of $M_D^2$. We infer that the lower limit on $\tan \beta$ coming from the estimated mass and mixing of the sterile neutrino dark matter is stronger in a wide range of parameter space compared to the bound coming from electroweak precision measurement computations [12] as shown in figure 4. For higher values of $M_D^2$ the upper limit on $\tan \beta$ becomes stronger than the upper bound coming from $\tau$ Yukawa coupling [12]. When combined these two limits provide a viable region for the Dirac wino mass.

Finally, it is important to check if this case [25] is consistent with the recent signal strengths, $\mu_{\gamma\gamma} = \frac{\sigma(pp \rightarrow h \rightarrow \gamma\gamma)}{\sigma(pp \rightarrow h \rightarrow \gamma\gamma)^{\text{SM}}}$ obtained by both CMS and ATLAS experiments or not. In this case the Higgs boson can decay to two light neutralinos namely the active and the sterile neutrino. The presence of light stops can provide certain enhancements in the signal strengths [25].

5.2. Case with large neutrino Yukawa coupling: $f \sim O(1)$
As previously noted for the case of the large Yukawa coupling we obtain a bino like light neutralino with mass around a few hundred MeV, which increases the total decay width of

2 We list the recent Higgs boson signal strengths in table 2
Figure 3. The red scattered points are obtained by varying the parameters in a wide range. All the shaded regions show different bounds obtained from different X-ray experiments [21, 27]. Finally the black star represents the central value of the observed mass and mixing by the XMM-Newton X-ray space observatory.

Figure 4. Showing the bound on $\tan \beta$ as a function of $M_2^D$.

the Higgs boson. The recent results from the CMS experiment shows a limit on the total decay width of Higgs as $k_{TOT} = \frac{\Gamma^{\mu \nu}_{R \rightarrow \mu \nu}}{\Gamma_{R \rightarrow \mu \nu}} \leq 5.4$ [26]. This result could constrains our parameter space severely [25].
| Channel         | $\mu$ (CMS)       | $\mu$ (ATLAS)       |
|-----------------|-------------------|---------------------|
| $h \rightarrow \gamma\gamma$ | $1.14^{+0.26}_{-0.23}$ [3] | $1.17^{+0.27}_{-0.27}$ [4] |
| $h \rightarrow ZZ^* \rightarrow 4l$ | $0.93^{+0.39}_{-0.32}$ [5] | $1.44^{+0.40}_{-0.33}$ [4] |
| $h \rightarrow WW^* \rightarrow 2l2\nu$ | $0.72^{+0.20}_{-0.18}$ [6] | $1.0^{+0.30}_{-0.20}$ [7] |
| $h \rightarrow bb$ | $1.0^{+0.27}_{-0.25}$ [8] | $0.2^{+0.60}_{-0.70}$ [9] |
| $h \rightarrow \tau\tau$ | $0.78^{+0.25}_{-0.22}$ [10] | $1.4^{+0.27}_{-0.4}$ [11] |

Table 2. Signal strengths ($\mu$) in different decay final states of the SM-like Higgs boson as reported by the CMS and the ATLAS collaborations (with the corresponding references).

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