Neutrino mass and proton decay in a $U(1)_R$-symmetric model

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We study a $U(1)_R$-symmetric extension of the supersymmetric standard model with supersymmetry breaking in the visible as well as hidden sectors. Specifically, we study $U(1)_R$-breaking effects parameterized by the gravitino mass. A special $R$-charge assignment of right-handed neutrinos allows us to have neutrino Yukawa couplings with the $R$-charged Higgs field, which develops a tiny vacuum expectation value after the inclusion of $U(1)_R$-symmetry breaking. Even with $O(1)$ Yukawa couplings, a suitable size of Dirac neutrino mass can be generated if the gravitino mass is very small, $m_{3/2} = 1–10$ eV. Our flipped $R$-charge assignment also allows a new type of dimension-five operator that can induce proton decay. It turns out that the proton stability mildly constrains the allowed range of the gravitino mass: A gravitino heavier than 10 keV can evade the proton decay constraint as well as cosmological ones. In this case, the largest neutrino Yukawa coupling is comparable to the electron Yukawa. We also calculate the mass of the pseudo goldstino and its mixing with neutralinos, and briefly discuss its implications in cosmology and Higgs phenomenology.

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1. Introduction

Weak-scale supersymmetry (SUSY) is an attractive candidate for the physics behind electroweak symmetry breaking (EWSB). SUSY breaking of the order of the weak scale can trigger the desired electroweak symmetry breaking either at tree level or via radiative corrections. However, null results at LHC for the SUSY particle search so far require us to reconsider such a picture, providing the motivation to extend the minimal SUSY standard model (MSSM).

The $U(1)_R$-symmetric extension of the SUSY standard model is an interesting starting point for the physics beyond the MSSM. A pleasing feature of $U(1)_R$ symmetry in the matter sector is that it naturally explains the absence of baryon-number-violating operators that would lead to fast proton decay; see Ref. [1] for a solution to the SUSY flavor problem. The gauge sector can be made $U(1)_R$ symmetric if each gaugino has Dirac mass, instead of the usual Majorana one. Such a Dirac mass term can be generated from hidden-sector SUSY breaking through a supersoft operator [2], which induces finite soft scalar masses. Models with the Dirac gaugino have the advantage of reducing the degree of fine tuning in the Higgs potential, even when the colored sparticles are as heavy as multi

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The Higgs sector can also be made $U(1)_R$ symmetric if we introduce the mirror partner, the $R$-partner of the MSSM Higgs doublets.

The origin of the EWSB can be addressed in a $U(1)_R$-symmetric manner. Reference [4] proposed\footnote{One of the motivations behind the construction in Ref. [4] was to explore the possibility of testing the SUSY breaking mechanism through the Higgs sector as a portal. Here we are interested in a simple realization of EWSB via the visible SUSY breaking.} a natural realization of the EWSB by coupling the $U(1)_R$-symmetric Higgs sector to visible-sector SUSY breaking, in which the supertrace sum rule is avoided by the presence of the hidden-sector SUSY breaking. Such coupling can be used to raise the lightest Higgs mass. The latter point is important if Dirac gaugino mass terms are generated by the supersoft operators, which also suppress the tree level $D$-terms, as was noted in Ref. [2]; see also Ref. [5] and references therein. The implications for cosmology as well as the Higgs search were also discussed in Ref. [6], in which it was pointed out that the visible SUSY breaking with Majorana gauginos is cosmologically disfavored.

In the present paper, we are interested in yet another notable aspect of the $R$-symmetric extension of the SUSY standard model. The $U(1)_R$ symmetry should be broken in the hidden sector for the cosmological constant to be canceled. If the $R$-symmetry breaking is mediated to the visible sector in a minimal way, an $R$-partner Higgs field develops a tiny vacuum expectation value (VEV) characterized by the gravitino mass. We will show that such a tiny VEV can be related to the smallness of neutrino masses. The generation of neutrino mass from SUSY breaking effects has been studied before in Refs. [7,8]. The generation of (Majorana) neutrino mass from $R$-symmetry breaking was also discussed recently in Refs. [9,10]. In the present paper, we will relate the tiny $R$-breaking VEV to Dirac neutrino masses by flipping the $R$-charge assignment of the right-handed neutrinos. Our flipped assignment of $U(1)_R$ charges allows us to write the desired coupling of the right-handed neutrino to the $R$-partner Higgs doublet; at the same time, it also allows a new type of dimension-five operator that makes the proton unstable. It turns out that the constraint from the proton stability is milder than the corresponding constraints in the MSSM. Thus, in our flipped $U(1)_R$ model, tiny neutrino masses can be generated without spoiling the feature of $U(1)_R$ symmetry of ensuring proton stability.

The present setup of the $U(1)_R$-symmetric model has potentially interesting implications for Higgs phenomenology. Since we consider the model of visible SUSY breaking and also assume that Dirac gaugino masses are induced from another SUSY breaking in a hidden sector, there appears a physical pseudo goldstino state $\xi$ [11], which directly couples to the Higgs sector. As discussed in Ref. [6], such pseudo goldstinos can affect the Higgs decay modes. We perform a similar analysis in our $U(1)_R$-symmetric setup and examine in some detail the mass and the mixing of the pseudo goldstino. It turns out that, although the pseudo goldstino can obtain mass due to the $R$-breaking effect, its mass $m_\xi$ is quite suppressed compared to the gravitino mass $m_{3/2}$. This is contrasted to the “sequestered” case in which $m_\xi$ is twice as large as $m_{3/2}$ [11,12]. Moreover, the mixings of the pseudo goldstino to the MSSM Higgsinos and gauginos are highly suppressed due to the softly broken $U(1)_R$ symmetry.

Another important constraint on the present model comes from cosmology; the light gravitino is constrained rather severely not to disturb successful big-bang nucleosynthesis (BBN), the structure formation of galaxies, and the cosmic microwave background (CMB) radiation, and not to be produced too much in the early universe. As we shall see later, the allowed range of the
gravitino mass is given by \[13–18\]

\[ m_{3/2} \lesssim 16 \text{ eV}, \quad 10 \text{ keV} \lesssim m_3/2. \]  

We will take account of these limits in the following analysis. We will also make a brief comment on the limit on the decaying dark matter scenario from the diffused gamma-ray line search.

The paper is organized as follows. In the next section, we present our model first in the \(R\)-symmetric limit, and then include \(R\)-breaking effects in a minimal way that are parameterized by the gravitino mass. Throughout the present paper, we assume the gravitino mass to be much smaller than the weak scale, although we do not elucidate possible origins of hidden-sector SUSY breaking. In Sect. 3, we show how the small VEVs of the \(R\)-charged Higgses can be used to explain the smallness of the neutrino mass. We also discuss the constraint from the proton decay in our flipped \(U(1)_R\) model for neutrino mass. In Sect. 4, we calculate the pseudo goldstino mass and the mixings to neutralinos under the assumption of minimal \(U(1)_R\) breaking. Specifically, we show how the mass of the pseudo goldstino is suppressed compared with the gravitino in the combined model of the Dirac gaugino and visible SUSY breaking. We then apply the results to briefly discuss the implications of the present model for Higgs phenomenology and cosmology. The final section is devoted to our conclusion and discussion. We also add two appendices concerning the analysis of baryon- and/or lepton-number-violating operators and a possible extension of the model in which the constraint from proton stability will be relaxed via \(U(1)\) flavor symmetry.

For simplicity, we will often refer to \(U(1)_R\) symmetry as “\(R\) symmetry”, if no confusion is expected.

2. \(R\)-symmetric model with visible SUSY breaking

In this section, we present an extension of the MSSM, in which \(U(1)_R\) symmetry is realized by combining the model of Dirac gauginos [2] with the model of visible SUSY breaking of Ref. [4]. We then discuss the effects of the minimal \(R\)-symmetry breaking induced from the coupling to supergravity.

2.1. \(U(1)_R\)-symmetric model

The \(U(1)_R\) symmetry forbids the usual Majorana gaugino mass terms as well as the Higgsino mass term, \(\mu H_u H_d\), in the MSSM. To realize \(U(1)_R\) symmetry, we extend the MSSM by introducing the \(R\)-partners of the gauginos and those of the Higgsinos.

First let us briefly discuss the gauge sector. Let \(a = 3, 2, 1\) parameterize each gauge group of \(SU(3)_C \times SU(2)_L \times U(1)_Y\). For each gauge group \(G_a\), we introduce an adjoint chiral multiplet \(A_a\), which contains an \(R\)-partner \(\chi_a\) of the gaugino \(\lambda_a\). Since the gaugino \(\lambda_a\) has \(R\)-charge \(+1\), its partner \(\chi_a\) has to have \(R\)-charge \(-1\) so that the adjoint chiral superfield \(A_a\) has \(R = 0\). Accordingly, we have a Dirac-type mass term

\[ \mathcal{L}_{\text{gaugino}} = - \sum_{a=1,2,3} m_a \lambda_a \chi_a + \text{H.c.} \]  

The Dirac gaugino mass term can be generated through the “supersoft” operator [2]

\[ \mathcal{L}_{\text{supersoft}} = \sum_{a=1,2,3} \int d^2 \theta \sqrt{2} \frac{W^a \bar{W}^a A_a}{\Lambda_D} + \text{H.c.}, \]

\(^2\)See Ref. [19] for further references. See also Sect. 4.3 for a brief discussion.
The absence of such a term can be justified if $U$ with the general argument $4$ and $U$ $\mu$ $/\Lambda_1$ where $\Lambda_1$ is the messenger scale, $W_a^g$ is the MSSM gauge field strength, $W'_a$ is the hidden-sector $U(1)'$ gauge field strength, which acquires a nonzero $D$-term, $\langle W'_a \rangle = \theta_a \langle D' \rangle$, so that $m_a = \langle D' \rangle / \Lambda_1$. In the present work, we assume that a suitable mass size is generated, although we do not elucidate the hidden-sector dynamics$^3$.

Next, we turn to the Higgs sector. Following Ref. [4], we consider the superpotential

$$W_{\text{Higgs}} = X_0(f + \lambda H_u H_d) + \mu_1 X_d H_u + \mu_2 X_u H_d.$$  \hspace{1cm} (4)

The gauge and $R$-charge assignments are shown in Table 1. The superfields $X_d$ and $X_u$ with $R$-charge 2 are the mirror partners of the MSSM Higgs fields $H_u$ and $H_d$, with the supersymmetric masses $\mu_1$ and $\mu_2$, respectively. The $SU(2) \times U(1)$ singlet field $X_0$ also has $R$-charge 2. The dimension-two parameter $f$ is the source of the visible SUSY breaking, while the dimensionless coupling $\lambda$ plays an important role not only in triggering the EWSB but also in generating the quartic coupling of the MSSM Higgs scalars.

With the $R$-symmetric superpotential (4), supersymmetry is spontaneously broken in accordance with the general argument$^4$. Moreover, such visible SUSY breaking triggers the correct EWSB if the coupling $\lambda$ is sufficiently large. A possible origin of the dimensionful parameters $f$ and $\mu_{1,2}$ was suggested in Ref. [4]. Here we just assume that the scale of these parameters is around the weak scale; we will take $\mu_{u,d} = 300$ GeV in our analysis in the next section.

We note that the adjoint chiral multiplets $A_{{\alpha}=2,1}$ of $SU(2) \times U(1)$ can have $U(1)_R$-symmetric superpotential interactions with the Higgs fields $H_{u,d}$ and $X_{d,u}$. For simplicity, however, we do not include them in our analysis. More importantly, we do not include the following mixing term between the singlet $X_0$ and the $U(1)_Y$ “adjoint” multiplet $A_1$:

$$W_{\text{mix}} = \mu_0 X_0 A_1.$$  \hspace{1cm} (5)

This term is very dangerous since it could cancel, if present, the linear term in the superpotential (4). The absence of such a term can be justified if $U(1)_Y$ is embedded into a simple gauge group at high energy. See Ref. [2] for a similar discussion about the absence of kinetic mixing of the hidden $U(1)'$ and $U(1)_Y$.

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$^3$ Dynamical SUSY breaking that realizes a non-vanishing $D$-term was recently discussed in Refs. [20,21].

$^4$ There exists no supersymmetric vacuum when the number of fields with $R$-charge 2 is larger than the number of fields with $R$-charge 0. See Ref. [22] and references therein.

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**Table 1.** The charge assignment of the Higgs and lepton fields under EW symmetry and $U(1)_R$ symmetry. The $R$-charges are those for left-handed chiral superfields. The MSSM Higgs doublets are $R$-neutral while their $R$-partners and the singlet $X_0$ have $R$-charge +2. All the quarks and leptons have $R$-charge +1, except that the right-handed neutrinos have $R$-charge $-1$. The implications of this flipped assignment will be discussed in Sect. 3. Note also that the adjoint chiral multiplets $A_\alpha$ are $R$-neutral so that their fermionic components have $R$-charge $-1$.

| | $X_0$ | $X_u$ | $X_d$ | $H_u$ | $H_d$ | $L$ | $E$ | $N$ |
|---|---|---|---|---|---|---|---|---|
| $SU(2)_L$ | 1 | 2 | 2 | 2 | 2 | 2 | 1 | 1 |
| $U(1)_Y$ | 0 | $+1/2$ | $-1/2$ | $+1/2$ | $-1/2$ | $-1/2$ | +1 | 0 |
| $U(1)_R$ | +2 | +2 | +2 | 0 | 0 | +1 | +1 | $-1$ |
2.2. Minimal $U(1)_R$-symmetry breaking

Next we discuss $U(1)_R$-symmetry breaking. Once the model is coupled to supergravity, the $U(1)_R$ symmetry is necessarily broken in order that the cosmological constant can be adjusted to zero. Our basic assumption here is that such breaking of $R$ symmetry is mediated to the visible sector in a minimal way; under such an assumption of “minimal $R$-breaking mediation”, the interaction Lagrangian is given by

$$\mathcal{L}_{\text{eff}} = \int d^2 \theta \phi^3 W + \text{H.c.}, \quad \phi = 1 + \theta^2 m_{3/2},$$

where $\phi$ is the so-called conformal compensator and $m_{3/2}$ is the mass of the gravitino. The conformal compensator $\phi$ can be absorbed by rescaling the chiral superfields, $\Phi_i \rightarrow \Phi_i$, where $\Phi_i$ represents all the chiral superfields in the theory. Classically, such rescaling has no effect on cubic terms in the superpotential, whereas terms of dimensions less than four are affected. In the present case, the interaction Lagrangian takes the form

$$\mathcal{L}_{\text{eff}} = \int d^2 \theta W + m_{3/2} G(\Phi_i) + \text{H.c.},$$

$$G(\Phi_i) \equiv 2 f X_0 + \mu_1 X_d H_u + \mu_2 X_u H_d,$$

where the fields in $G(\Phi_i)$ are the scalar components of the corresponding superfields $\Phi_i = H_u$, $H_d$, $X_u$, $X_d$, or $X_0$. Terms proportional to the gravitino mass represent $R$-breaking interactions, and are small if the gravitino mass is small; we assume throughout the present paper that the gravitino mass is much smaller than the weak scale. As we will see shortly, however, these interactions induce a slight shift of the vacuum since they contain a tadpole term of $X_0$ and also those of $X_{u,d}$ after the EWSB. Such a VEV shift will play an important role when we discuss the generation of Dirac neutrino masses (in Sect. 3) and the properties of the pseudo goldstino (in Sect. 4).

2.3. The EWSB vacuum

We now analyze the scalar potential in order to find the shift of the VEVs induced by the $R$-breaking effects. Here we assume that the Kähler potential is canonical. The scalar potential involving the neutral Higgs fields $H_{u,d}^0$, their mirror partners $X_{u,d}^0$, and the singlet $X_0$ is given by

$$V(\Phi_i, \Phi_i^\dagger) = V_0(\Phi_i, \Phi_i^\dagger) - \{ m_{3/2} G(\Phi_i) + \text{H.c.} \},$$

with the $R$-symmetric part $V_0$ given by

$$V_0 = | f - \lambda H_{d}^0 H_{d}^0 |^2 + | \lambda X_0 H_{d}^0 - \mu_1 X_d H_{d}^0 |^2 + | \lambda X_0 H_u^0 + \mu_2 X_u H_d^0 |^2$$

$$+ \frac{1}{8} \theta^2 (| H_{u}^0 |^2 - | H_{d}^0 |^2 | X_{u}^0 |^2 - | X_{d}^0 |^2 | X_d^0 |^2)$$

$$+ \tilde{m}_d^2 | H_{u}^0 |^2 + \tilde{m}_d^2 | H_{d}^0 |^2 - (b H_{d}^0 H_{d}^0 + h.c.)$$

$$+ m_{X_0}^2 | X_0 |^2 + m_{X_d}^2 | X_d^0 |^2 + m_{X_u}^2 | X_u^0 |^2.$$

Here we assume that the chiral compensator takes the conventional form whose $F$-component $F_\phi$ is equal to the gravitino mass after solving the equation of motion for $F_\phi$. We add that this is not always true, for instance, as in the models with “almost no-scale” SUSY breaking [23], where linear terms in the Kähler potential play an important role. Here we assume that our hidden-sector SUSY breaking does not involve such linear terms. See also Refs. [24,25], which proposed an improved gauge fixing of conformal supergravity and discussed the conditions under which the conventional expression (6) can be justified.
The first and second lines represent the $F$-term and $D$-term potentials, respectively. The coupling constant $g^2$ is $g_1^2 + g_2^2$, where $g_1$ and $g_2$ are the gauge coupling constants of $U(1)$ and $SU(2)$ symmetries, respectively. The remaining are soft SUSY breaking terms, except that the masses of the MSSM Higgs fields, $m_{H_\pm}^2$, are the sums of the soft masses and the supersymmetric masses. We suppose that these soft terms are induced radiatively from the supersoft operators, or mediated from the hidden sector in an $R$-invariant way. We also assume that all dimensionful parameters are around or just above the weak scale, whereas the gravitino mass is much smaller. In the following, we consider the case in which $m_{X_0}^2$ and $m_{X_u,d}^2$ are positive so that $X_0$ and $X_{u,d}$ do not develop VEVs in the $R$-symmetric limit.

The vacuum with EWSB can be found by solving the stationary conditions

$$0 = \frac{\partial V}{\partial \Phi_i} = \frac{\partial V_0}{\partial \Phi_i} - m_{3/2}^2 \frac{\partial G(\Phi_i)}{\partial \Phi_i}. \quad (11)$$

We solve these equations perturbatively with respect to a small gravitino mass. In the $R$-symmetric limit, i.e., $m_{3/2} \rightarrow 0$, the solution was given in Ref. [4], which we denote by $\langle \Phi_i \rangle_0$. Note that the unbroken $U(1)_R$ symmetry implies $\langle X_0 \rangle_0 = \langle X_u \rangle_0 = \langle X_d \rangle_0 = 0$.

Now, with the $R$-breaking terms, the solution to Eq. (11) takes the form

$$\langle \Phi_i \rangle = \langle \Phi_i \rangle_0 + \kappa_i m_{3/2} + O(m_{3/2}^2), \quad (12)$$

where the coefficient $\kappa_i$ is found to be

$$\kappa_i = M^{-1}_{ij} \frac{\partial G}{\partial \Phi_j} \bigg|_{\Phi = \langle \Phi \rangle_0}, \quad M_{ij} \equiv \frac{\partial^2 V_0}{\partial \Phi_i \partial \Phi_j} \bigg|_{\Phi = \langle \Phi \rangle_0}. \quad (13)$$

Explicitly, we obtain, for the $R$-charged Higgs fields,

$$\langle X_0 \rangle = \frac{2f}{m_{X_0}^2 + \lambda^2 v^2 m_{3/2}^2}, \quad (14)$$

$$\langle X_0^2 \rangle = \frac{\mu_1 v \sin \beta}{m_{X_d}^2 + \mu_1^2 + (m_Z^2/2) \cos 2\beta} m_{3/2}^2, \quad (15)$$

$$\langle X_0^2 \rangle = \frac{-\mu_2 v \cos \beta}{m_{X_u}^2 + \mu_2^2 - (m_Z^2/2) \cos 2\beta} m_{3/2}^2, \quad (16)$$

where $m_Z^2 = (g_1^2 + g_2^2)v^2/2$ is the $Z$ boson mass. We see that the VEV of $X_0$ is proportional to $f$, which represents the scale of the visible SUSY breaking, while the VEVs of $X_{u,d}$ are proportional to their supersymmetric masses and their partner’s VEVs. These proportionality can be understood if one notices that replacing the Higgs fields by their VEVs in Eq. (8) generates the tadpole term for each of $X_{0,u,d}$:

$$G(X_0, X_{u}^0, X_{d}^0) = 2f X_0 + \mu_1 \langle H_u^0 \rangle X_{d}^0 - \mu_2 \langle H_d^0 \rangle X_{u}^0. \quad (17)$$

As noted in Ref. [2], the $SU(2) \times U(1)$ $D$-terms are absent in the supersoft limit. If the adjoint scalars get soft scalar masses through $R$-invariant mediation of SUSY breaking, then the $D$-terms do not decouple completely. Our results in the present paper are not affected whether the $D$-terms decouple or not.

Note that the VEVs of the MSSM Higgs fields also receive $O(m_{3/2})$ shifts, which are negligibly small for a small gravitino mass.
Of course, the VEVs of these \( R \)-charged fields are all proportional to the gravitino mass\(^8\), which parameterizes the \( U(1)_R \)-symmetry breaking.

3. Neutrino mass and proton decay in the flipped \( U(1)_R \) model

In this section, we show that neutrinos can acquire tiny masses via the supergravity-induced effect of \( U(1)_R \)-symmetry breaking. We introduce three right-handed neutrino multiplets \( N^c_i \) \((i = 1, 2, 3)\) and assign the \( R \)-charge \(-1\) to them, as shown in Table 1. That is, we flip the \( R \)-charge of the right-handed neutrinos. Under this flipped assignment, we first argue that the neutrinos can be Dirac particles, and that the smallness of the neutrino masses is related to the smallness of \( U(1)_R \)-symmetry breaking, parameterized by the gravitino mass. Then, we discuss the proton decay induced by a new type of dimension-five operator involving the right-handed neutrinos.

3.1. Neutrino mass in the flipped \( U(1)_R \) model

Under our flipped \( U(1)_R \) assignment, the right-handed neutrino can have a Yukawa-type interaction with the mirror Higgs field \( X_u \), instead of the MSSM Higgs \( H_u \). At the renormalizable level, the matter superpotential is given by

\[
W = y_U^{ij} Q_i U^c j H_u + y_D^{ij} Q_i D^c j H_d + y_E^{ij} E^c i L_j H_d + y_N^{ij} N^c i L_j X_u ,
\]

where \( y_{U,D,E,N} \) are \( 3 \times 3 \) Yukawa coupling matrices.

The Majorana mass term for the right-handed neutrinos is forbidden by our \( R \)-charge assignment. The symmetry allows the Weinberg operator

\[
W_{\Delta L=2} = - \frac{c^{ij}}{\Lambda_{\text{cutoff}}} (L_i H_u)(L_j H_u) ,
\]

which makes a sub-dominant contribution to the neutrino mass matrix if we assume the cutoff \( \Lambda_{\text{cutoff}} \) to be larger than \( 10^{12} - 10^{13} \) GeV. Therefore the neutrinos are almost Dirac particles in our model.

By replacing the mirror Higgs field with its VEV (16), we obtain the Dirac mass term for neutrinos as

\[
L_{\nu \text{ mass}} = -m_\nu v^c_R v_L + \text{H.c.}, \quad m^{ij}_\nu = y_N^{ij} (X_u^0) .
\]

Since we are most interested in the heaviest neutrino, we will suppress the flavor indices hereafter and denote the heaviest eigenvalue and the corresponding Yukawa coupling by \( m_\nu \) and \( y_\nu \), respectively. With this understanding, we write

\[
m_\nu = y_\nu (X_u^0) = y_\nu m_{3/2} \left[ \frac{-\mu_2 v \cos \beta}{m_X^2 + \mu_2^2 - (m_Z^2/2) \cos 2\beta} \right] .
\]

A notable feature of the present model is that the size of the neutrino mass is set by that of the gravitino mass \( m_{3/2} \) when all other dimensionful parameters are of the same order. Accordingly, the neutrino Yukawa coupling can be of \( O(1) \) if the gravitino is as light as 10 eV. In fact, as we mentioned

\(^8\) Precisely speaking, the VEVs \( (X_{0,u,d}) \) have \( R \)-charge +2 and should be proportional to \( m_{3/2}^* \) since one can regard the gravitino mass as having \( R \)-charge \(-2\) in the sense of spurion analysis.
Fig. 1. The heaviest neutrino mass as a function of the Higgsino mass parameter $\mu_2$, with $m_{X_d} = 250$ GeV (left) and 500 GeV (right). In each figure, the solid and dotted lines correspond to the case with and without the $D$-term contribution in the denominator in Eq. (21). The three lines correspond to $\tan \beta = 3$ (top), $\tan \beta = 5$ (middle), and $\tan \beta = 10$ (bottom), respectively.

In Eq. (1), the mass of the gravitino is constrained cosmologically:

$$m_{3/2} \lesssim 16 \text{ eV}, \quad 10 \text{ keV} \lesssim m_{3/2}.$$ 

In the case of a lighter gravitino, a tiny neutrino mass can be obtained with a large Yukawa coupling. Even in the case of a heavier gravitino, the neutrino Yukawa can be comparable to the electron Yukawa.

Figure 1 shows how the generated neutrino mass depends on the Higgs mass parameters, for a fixed value of $m_{3/2} = 10$ eV and $y_\nu = 0.1$. We see that the neutrino mass is maximized for $\mu_2 \sim m_{X_d}$. We also note that the $D$-term contribution in the denominator in Eq. (21) can be safely neglected when the Higgs mass parameters are larger than $m_Z$. Accordingly, the neutrino mass is proportional to $\cos \beta$ instead of $\sin \beta$.

Some remarks are in order. First, our $U(1)_{R}$-charge assignment does not induce mixed anomalies with the SM gauge groups since the right-handed neutrinos are SM singlets. $U(1)_{R}$ anomalies can be canceled by introducing other singlets with positive $U(1)_{R}$-charges, without changing our results below. Second, our model can be regarded as a realization of the “neutrinophilic Higgs” idea: the original idea was proposed in a non-SUSY context in Refs. [26,27], in which a softly broken $Z_2$ symmetry is the source of the tiny VEV; see also Ref. [28] for a softly broken global $U(1)$ case. Quantum stability was discussed in Refs. [29,30]. The point in our model is that the size of the “neutrinophilic” Higgs VEV is related to that of the gravitino mass, and we can address its implications for cosmology and Higgs phenomenology, in addition to the proton stability, as we shall discuss shortly. We also note that many other topics have been discussed in the literature, such as muon $g - 2$ and lepton flavor violations [27,28], low-scale leptogenesis [31], dark matter and cosmology [32], dark energy [33], and supernova neutrinos [34]: In particular, Ref. [35] claimed that the sizes of neutrino Yukawa couplings are severely constrained from the observations of supernova neutrinos as well as CMB radiations, if the neutrinophilic Higgs scalar is extremely light. This constraint does not apply here since the corresponding scalar is heavy enough.

3.2. Proton decay in the flipped $U(1)_{R}$ model

Originally, one pleasing feature of $U(1)_{R}$ symmetry was that it explains proton stability naturally: the dangerous baryon- and lepton-number-violating operators are forbidden if we assign $R$-charge +1 to all the quark and lepton superfields and $R$-charge 0 to the Higgses. This property is not modified
when we introduce the $R$-partner Higgses $X_{0,u,d}$ with $R = 2$ and the $R$-partner gauginos with $R = 0$. In our assignment, however, we flip the $R$-charge of the right-handed neutrinos $N^c$ to be $-1$, which may spoil the proton stability.

An operator analysis presented in Appendix A shows that there is a unique dimension-five operator

$$W_{5R} = -\frac{C_{ijkl}^{5}{\Lambda}_{\text{cutoff}}}{U_{i}^{c}D_{j}^{c}D_{k}^{c}N_{\ell}^{c}},$$

(22)

where $\Lambda_{\text{cutoff}}$ is the cutoff scale and $C_{ijkl}^{5}$ is a dimensionless coefficient. In this subsection, we will give a rough estimate of the proton lifetime. To simplify the expressions, we will work in a flavor basis in which the Yukawa couplings of up-type quarks and charged leptons are diagonal; we also define $C_{ijkl}^{5}$ as the coefficient of the $\ell$th mass eigenstate of neutrinos.

With the dimension-five operator (22), the proton decay occurs via the processes depicted in Fig. 2, one with neutral Higgsino dressing and the other with charged Higgsino dressing. Remarkably, diagrams involving the top Yukawa coupling are absent because the $\mu H_u H_d$ term is absent in the $U(1)_R$-symmetric limit, or is extremely suppressed, as it is proportional to the tiny VEV (14). Instead, we have potentially large contributions involving a large neutrino Yukawa coupling $y_\nu$ of the heaviest neutrino $\nu_3$.

The partial rate for a proton decaying into a meson $M$ and a lepton $\ell$ takes the form

$$\Gamma(p \to M + \ell) = \frac{m_p}{32\pi} \left(1 - \frac{m_M^2}{m_p^2}\right)^2 \frac{|\alpha_p|^2}{f_M^2} |A(p \to M \ell)|^2.$$

(23)

Here $M_M$ and $f_M$ are the mass and the decay constant of the meson $M$, respectively. The dimensionful constant $\alpha_p$ is defined through $\langle 0 | \epsilon_{abc} (d_R^a u_R^b) u_L^c | 0 \rangle = \alpha_p N_L$ where $a, b, c$ are color indices and $N_L$ is the wavefunction of a left-handed proton. We use the value $\alpha_p = -0.015 \text{ GeV}^3$ [37] in our calculation. The amplitude $A$ corresponding to each diagram in Fig. 2 is given respectively by

$$\begin{align*}
A(p \to K^+ \nu_3) &= \frac{C_{1123}^{1123}}{\Lambda_{\text{cutoff}}} \frac{y_\nu y_3 \mu_2}{16\pi^2 m_{\text{soft}}^2}, \\
A(p \to \pi^0 c^+) &= \sum_{k=s,b} \frac{C_{5}^{11k3}}{\Lambda_{\text{cutoff}}} \frac{U_{\ell k} y_3 V_{ak} y_k \mu_2}{16\pi^2 m_{\text{soft}}^2},
\end{align*}$$

(24)

(25)

where $U_{\ell3}$ and $V_{ak}$ are mixing matrix elements, and $m_{\text{soft}}$ is a typical mass scale of the particles propagating in the loop. We see that the latter diagrams with charged Higgsino dressing are suppressed by

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9 For the long-distance effects of an $RRRR$-type operator, see Ref. [36].
the CKM matrix element $\varepsilon = V_{us} = 0.22$, because $y_t \sim y_b \varepsilon^2$ and $y_b V_{ub} \sim y_t V_{us} \sim y_b \varepsilon^3$. Therefore, the dominant decay mode is $p \rightarrow K^+ + \bar{\nu}$. Using the simplified expression of the neutrino Yukawa\textsuperscript{10}

$$y_\nu \approx \frac{m_\nu}{m_{3/2}} \left( \frac{m^2_{X_u} + \mu_2^2}{-\mu_2 v \cos \beta} \right),$$

we find that the partial decay width into a kaon and an anti-neutrino is given by

$$\Gamma(p \rightarrow K^+ \bar{\nu}) \approx \frac{m_p}{32\pi} \left(1 - \frac{m_{K^+}^2}{m_p^2}\right)^2 \frac{\alpha_p^2}{f_K^2} \left(\frac{C_5}{16\pi^2 m_{3/2} \Lambda_{\text{cutoff}}} \right) \left(\frac{m^2_\nu m^2_{X_u}}{v^4 \cos^4 \beta} \right) \left(\frac{\mu_2^2 + m^2_{X_u}}{m^2_{\text{soft}}} \right)^2,$$

(27)

where we denote $C_5^{1123}$ simply by $C_5$.

Numerically, the lifetime of a proton can be estimated as

$$\tau_p = \frac{5.65 \times 10^{28} \text{yr}}{C_5^2} \left(\frac{\Lambda_{\text{cutoff}}}{2.4 \times 10^{18} \text{GeV}}\right) \left(\frac{m_{3/2}}{10 \text{eV}}\right) \left(\frac{\cos^2 \beta}{0.1}\right)^2,$$

(28)

where we took the following parameter values:

$$m_{\text{soft}} = 1.5 \text{ TeV}, \quad \mu_2 = m_{X_u} = 300 \text{ GeV}, \quad \tan \beta = 3,$$

$$m_p = 1.0 \text{ GeV}, \quad \alpha_p = -0.015 \text{ GeV}^3,$$

$$m_{K^+} = 0.5 \text{ GeV}, \quad f_K = 0.13 \text{ GeV},$$

$$m_s = 0.1 \text{ GeV}, \quad m_\nu = 0.1 \text{ eV}.$$ (29)

The present lower bound on the proton lifetime is $2.3 \times 10^{33} \text{ yr} (90\% \text{ CL})$, obtained for the $p \rightarrow K^+ + \bar{\nu}$ mode\textsuperscript{38}. By comparing the bound with the result, we obtain the constraint\textsuperscript{11}

$$|C_5| \lesssim 4.96 \times 10^{-3} \left(\frac{\Lambda_{\text{cutoff}}}{2.4 \times 10^{18} \text{GeV}}\right) \left(\frac{m_{3/2}}{10 \text{eV}}\right) \left(\frac{\cos^2 \beta}{0.1}\right).$$

(30)

Figure 3 shows the constraint from proton stability as a function of the gravitino mass. The constraint is most severe for a small gravitino mass: If the gravitino is as light as $\mathcal{O}(\text{eV})$, the coefficient $C_5$ needs to be a few orders of magnitude smaller than 1 even when the cutoff $\Lambda_{\text{cutoff}}$ is equal to the reduced Planck scale $M_{\text{pl}}$. The constraint becomes milder for a larger gravitino mass. This can be understood from Eq. (21): For a fixed neutrino mass, a larger gravitino mass implies a smaller neutrino Yukawa coupling, suppressing the proton decay. We thus find that the constraint can be satisfied for a gravitino mass larger than 10 keV, which is also viable cosmologically, as in seen in Eq. (1).

Notice, however, that the constraint becomes severe for a large tan $\beta$: The lifetime is proportional to $\cos^4 \beta$ up to a possible mild dependence in the denominator in Eq. (21). This is because the $X^0_u$ tadpole is proportional to $\langle H^0_u \rangle$ and because the decay amplitude involves a down-type Yukawa coupling. Consequently, for $\tan \beta = 10$, the proton stability against the dimension-five operator (22) requires that the gravitino mass should be larger than $\mathcal{O}(100 \text{ keV})$. In this case, the neutrino Yukawa coupling $y_\nu$ is comparable to the electron Yukawa.

\textsuperscript{10} Here we have neglected the mild tan $\beta$ dependence in the denominator in Eq. (21) to simplify the expression. We used Eq. (21) in our numerical calculation.

\textsuperscript{11} This only applies for $\Lambda_{\text{cutoff}} \gtrsim 10^{13} \text{ GeV}$; otherwise, the operators (19) can no longer be neglected.
Fig. 3. Constraints from proton stability as a function of the gravitino mass. Left: the constraint on \( \log_{10} \Lambda_{\text{cutoff}} \) for the fixed coupling \( C_5 = 1 \). Right: the constraint on \( \log_{10} |C_5| \) for \( \Lambda_{\text{cutoff}} = M_{\text{pl}} \). We took \( \tan \beta = 3, \mu_2 = 300 \text{ GeV}, m_{X_{\nu}} = 300 \text{ GeV}, \) and \( m_{\tilde{q}} = 1.5 \text{ TeV} \). The figures correspond to the neutrino mass \( m_{\nu} = 0.1 \text{ eV} \). The gray regions are excluded by the proton decay while the yellow bands are disfavored cosmologically.

The above results should be compared with the MSSM case in which the proton decay is induced by the dimension-five operators

\[
\Delta W_{\text{MSSM}} = -\frac{C_L}{2 \Lambda_{\text{cutoff}}} Q Q L - \frac{C_R}{\Lambda_{\text{cutoff}}} U^c D^c E^c.
\]

In the MSSM, the dominant contribution comes from the \( L L L L \) operator, and the coefficient should be very suppressed, \( |C_L| \lesssim 10^{-8} \) for \( \Lambda_{\text{cutoff}} = M_{\text{pl}} \). Moreover, the \( R R R R \) operator should also be suppressed because it can involve the top Yukawa coupling when dressed with a charged Higgsino loop. In our case, the \( L L L L \) operator is absent due to the \( R \)-symmetry, and with our \( R R R R \) operator, chargino loop diagrams in Fig. 2 involve a down-type Yukawa coupling, \( y_s \) or \( y_b V_{cb} \), instead of top Yukawa. As a consequence, the \( R R R R \) contribution in the flipped \( U(1)_R \) model is a couple of orders smaller than that in the MSSM.

In this way, our flipped \( U(1)_R \) assignment can explain the smallness of neutrino masses with relatively large Yukawa couplings without spoiling the proton stability so much. The latter is certainly true for small \( \tan \beta \): for larger \( \tan \beta \), on the other hand, there is a tension between our mechanism for neutrino masses and the proton stability.

We note that such tension can be relaxed if we combine the idea of flavor symmetries with the present model. In Appendix A, we give an illustrative example of \( U(1) \) flavor symmetry, along the lines of Ref. [39], and show that a proper \( U(1) \) charge assignment for generating the Yukawa hierarchy can guarantee the proton stability as well.

4. Mass and mixing of the pseudo goldstino

We now turn to another effect of \( R \)-symmetry breaking in our \( U(1)_R \)-symmetric model. When supersymmetry is broken in two independent sectors, there appear two Goldstone fermions [11]. After coupling to supergravity, one linear combination of these goldstinos becomes the longitudinal components of the massive gravitino, and the other is a pseudo goldstino state, which we denote by \( \zeta \). The mass of such a pseudo goldstino has been studied in the literature [6,11,12].

In our case, supersymmetry is broken in the visible sector as well as in a hidden sector. For our purpose, we do not need to specify the precise form of the hidden sector; we just assume that the supersymmetry breaking is hierarchical: the SUSY breaking scale in the hidden sector is much larger
than that in the visible sector. We then expect that the physical pseudo goldstino state $\zeta$ resides dominantly in the visible sector. In the limit $\lambda \to 0$, in which the EWSB is switched off, the would-be goldstino in the visible sector is the singlino $\sim X_0$, the fermionic component of the singlet $X_0$, since the visible SUSY breaking is triggered by its linear term in the superpotential (4). Therefore we first discuss the mass of the $X_0$ fermion, and then examine the full neutralino mass matrix to calculate its smallest eigenvalue and the corresponding eigenvector.

4.1. Mass of the singlet fermion from $U(1)_{R}$ breaking

The mass term of the singlino $\sim X_0$ can be generated from the contact term in the Kähler potential of $X_0$. For definiteness, let us consider the following Kähler potential:

$$K(X_0, X_0^\dagger) = X_0 X_0^\dagger - \frac{1}{4\Lambda_0^2} (X_0 X_0^\dagger)^2,$$  (32)

where $\Lambda_0$ is the cutoff scale at which the contact term is generated. The relevant terms in the Lagrangian containing the $X_0$ supermultiplet are given by

$$\mathcal{L}_{X_0} = K_{X_0 X_0} F_{X_0} F_{X_0} - m_{X_0}^2 X_0 X_0^\dagger - \lambda^2 (|H_u|^2 + |H_d|^2) X_0 X_0^\dagger X_0 X_0^\dagger$$

$$+ \left\{ F_{X_0} \frac{\partial W}{\partial X_0} - \frac{1}{2} m_{\tilde{X}_0} \tilde{X}_0 \tilde{X}_0 + \text{H.c.} \right\} + \{2 f m_{3/2} X_0 + \text{H.c.} \} ,$$  (33)

where the (Majorana) mass of the would-be goldstino $\sim X_0$ is found to be

$$m_{\tilde{X}_0} = -\frac{1}{2\Lambda_0^2} \left\langle \int d^2 \theta \frac{1}{2} X_0 X_0^\dagger \right\rangle = -\frac{\langle F_{X_0}\rangle^\dagger}{2\Lambda_0^2} \langle X_0 \rangle.$$  (34)

Note that this expectation value can be non-vanishing only in the presence of the $U(1)_R$ breaking, the last term in the Lagrangian (33). To compute it, we use the equations of motion to get

$$\langle F_{X_0}\rangle = -K_{X_0 X_0}^{-1} \frac{\partial W}{\partial X_0} \approx - \left( f - \lambda \langle H_u^0 \rangle \langle H_d^0 \rangle \right) ,$$  (35)

$$\langle X_0 \rangle = \frac{2f}{m_{X_0}^2 + \lambda^2 v^2 + \delta m_{X_0}^2} m_{3/2} .$$  (36)

Here, in the first equation, we have used the fact that the Kähler metric is almost canonical since $|\langle X_0 \rangle| \ll \Lambda_0$. In the second equation, we note that the previous result (14) is slightly modified by the soft scalar mass $\delta m_{X_0}^2$ due to the contact term in Eq. (32):

$$\delta m_{X_0}^2 \equiv \frac{|\langle F_{X_0}\rangle|^2}{\Lambda_0^2} \approx \frac{1}{\Lambda_0^2} |f - \lambda v^2 \sin \beta \cos \beta|^2 .$$  (37)

By plugging the VEVs of $X_0$ and $F_{X_0}$ into expression (34), we thus find that the $X_0$ fermion mass is proportional to the gravitino mass as

$$m_{\tilde{X}_0} \approx 2 \mathcal{Z}_1 m_{3/2}, \quad \mathcal{Z}_1 \equiv \frac{f (f - \lambda v^2 \sin \beta \cos \beta)}{\Lambda_0^2 m_{X_0}^2 + \lambda^2 v^2 + \delta m_{X_0}^2} .$$  (38)

The result (38) is consistent with the general assertion [11] that the pseudo goldstino mass is twice the gravitino mass in the “sequestered” limit. This can be seen as follows: If the soft mass of $X_0$ is
zero, $m_{X_0}^2 \to 0$, and if the Higgs sector decouples from the visible SUSY breaking, $\lambda \to 0$, we have $\delta m_{X_0}^2 \to f^2/\Lambda_0^2$, and hence,

$$Z_1 \to 1.$$  \hfill (39)

Otherwise, the pseudo goldstino mass is suppressed by a factor $Z_1$, which is roughly of order $f/\Lambda_0^2$, in accordance with the general argument [6]. For instance, for a moderate choice of parameters $f \sim (10^3 \text{ GeV})^2$ and $\Lambda_0 \sim 10^7 \text{ GeV}$, the mass of the pseudo goldstino is smaller than $O(10 \text{ eV})$ even if the gravitino is as heavy as $m_{3/2} \sim \text{ GeV}$.

4.2. Neutralino mass matrix with the Dirac gaugino

Now we consider the neutralino mass matrix that incorporates the pseudo goldstino. In this subsection, let us denote by $\chi_B$ and $\chi_W$ the Dirac partners of the bino $\tilde{B}$ and the wino $\tilde{W}$, respectively. The Dirac mass terms for $SU(2) \times U(1)$ gauginos are

$$\mathcal{L}_{\text{gaugino}} = -m_{\tilde{B}} \chi_B \chi_B - m_{\tilde{W}} \chi_W \chi_W + H.c..$$  \hfill (40)

Then, in a basis given by

$$\tilde{\Psi}^T = (\chi_B \chi_W \tilde{W}^3 \tilde{B} | \tilde{H}^0_d \tilde{H}^0_u \tilde{X}^0_d \tilde{X}^0_u | \tilde{X}_0),$$  \hfill (41)

the neutralino mass matrix takes the form

$$\mathcal{M}_{\tilde{N}} = \begin{pmatrix} 0 & 0 & m_{\tilde{B}} & 0 & 0 & 0 & 0 \\ 0 & m_{\tilde{W}} & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & am_Z & bm_Z & a'm_{V} & b'm_{V} & 0 \\ 0 & m_{\tilde{W}} & 0 & cm_Z & dm_Z & c'm_{V} & d'm_{V} \\ 0 & 0 & am_Z & cm_Z & 0 & -\lambda v_0 & 0 & +\mu_2 & -\lambda v_u \\ 0 & 0 & bm_Z & dm_Z & -\lambda v_0 & 0 & -\mu_1 & 0 & -\lambda v_d \\ 0 & 0 & a'm_{V} & c'm_{V} & 0 & -\mu_1 & 0 & 0 & 0 \\ 0 & 0 & b'm_{V} & d'm_{V} & +\mu_2 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & -\lambda v_u & -\lambda v_d & 0 & 0 & m_{\tilde{X}_0} \end{pmatrix},$$  \hfill (42)

where $m_{\tilde{X}_0} = 2m_{3/2}Z_1$, and we have defined $\langle X^0 \rangle = v_0$ and

$$m_{V}^2 = \frac{g_1^2 + g_2^2}{2} (|\langle X^0_d \rangle|^2 + |\langle X^0_u \rangle|^2), \quad \tan \gamma = \frac{\langle X^0_u \rangle}{\langle X^0_d \rangle}. $$  \hfill (43)

We have also used the following abbreviations, with $\theta_W$ being the weak mixing angle:

$$U = \begin{pmatrix} a & b \\ c & d \end{pmatrix} = \begin{pmatrix} c_\beta c_W & -s_\beta c_W \\ -c_\beta s_W & s_\beta s_W \end{pmatrix} = \begin{pmatrix} \cos \theta_W & -\sin \theta_W \\ \sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} \cos \beta & -\sin \beta \end{pmatrix},$$  \hfill (44)

where $\langle H^0_d \rangle = v \cos \beta$ and $\langle H^0_u \rangle = v \sin \beta$; the primed quantities are obtained by replacing $\beta$ with the angle $\gamma$ defined above.

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The mass eigenvalue $\Omega_D$ can be found by solving the characteristic equation

$$0 = \det(M_N - \Omega_D \hat{1}) = \det M_N - \Omega_D \det M_N \text{Tr}[M_N^{-1}] + \cdots,$$

(45)

where the dots represents terms of a higher order in $\Omega_D$. The first and second terms are calculated to be

$$\det M_N = 2Z_1 m_{3/2} \mu_1^2 \mu_2^2 \omega_B^2 \omega_W^2,$$  

(46)

$$\det M_N \text{Tr}[M_N^{-1}] = m_{3/2} \omega_B^2 [\mu_1^2 \mu_2^2 + \lambda^2 v^2 (\mu_1^2 \sin^2 \beta + \mu_2^2 \cos^2 \beta)],$$

(47)

up to $O(m_{3/2}^2)$ and $O(m_{3/2})$, respectively. The lowest eigenvalue, which we identify with the pseudo goldstino mass $m_\zeta$, can be calculated by keeping only the first two terms in Eq. (45),

$$m_\zeta = 2m_{3/2} Z_1 Z_2,$$

(48)

where the suppression factor $Z_1$ is given in Eq. (38) and

$$Z_2^{-1} = 1 + \frac{\lambda^2 v^2 \sin^2 \beta}{\mu_2^2} + \frac{\lambda^2 v^2 \cos^2 \beta}{\mu_1^2}.$$  

(49)

We see that the mass of the true pseudo goldstino is suppressed by a factor $Z_1 Z_2$. We note that the lightness of the pseudo goldstino will be protected against quantum corrections by the $U(1)_R$ symmetry in the visible sector, as long as the Dirac nature of the gauginos is kept.

The corresponding eigenvector can be found perturbatively in $R$-breaking effects. By writing $\bar{\Psi}_\zeta = \bar{\Psi}_0 + \bar{\Psi}_1 + \cdots$, we find

$$\bar{\Psi}_0^T = \sqrt{Z_2} \begin{pmatrix} 0 & 0 & 0 & 0 & 0 \end{pmatrix} (\zeta^{(0)}_X \zeta^{(0)}_H) \frac{1}{Z_1},$$

$$\bar{\Psi}_1^T = \sqrt{Z_2} (\zeta^{(1)}_X \zeta^{(1)}_H) \begin{pmatrix} 0 & 0 & 0 \end{pmatrix}.$$  

(50)

The two-component vectors $\zeta^{(0)}_X, \zeta^{(1)}_X$, and $\zeta^{(1)}_H$ are given respectively by

$$\zeta^{(0)}_X = \begin{pmatrix} -\frac{\lambda v \mu_1}{\mu_2} \\ \frac{m_B}{m_B} \end{pmatrix}, \quad \zeta^{(1)}_X = \frac{m_{3/2}}{\cos \theta_W} \begin{pmatrix} \sin \theta_W \\ \frac{m_B}{m_B} \end{pmatrix}, \quad \zeta^{(1)}_H = \frac{m_\zeta}{\frac{\lambda v \mu_1}{\mu_2}} \begin{pmatrix} \frac{\lambda v \mu_1}{\mu_2} \\ \frac{m_B}{m_B} \end{pmatrix},$$

(51)

where the dimensionless factor $Z_3$ is defined by

$$Z_3 = -m_Z s_\beta c_\beta \left( \frac{\lambda v}{\mu_2} - \frac{\lambda v}{\mu_1^2} m_{3/2} \right) + \lambda v \left( \frac{s_\beta s_\gamma}{\mu_2} + \frac{c_\beta c_\gamma}{\mu_1} \right) \frac{m_V}{m_{3/2}}.$$  

(52)

We see that, at the leading order, only the xinos $\tilde{X}_{d,u}^0$ can mix with the would-be goldstino $\tilde{X}_0$ because of the $U(1)_R$ charge conservation; at the next order, the $\tilde{X}_0$ mixes with the fermions with $R = -1$, i.e., the Higgsinos $\tilde{H}_{d,u}^0$ and the gaugino Dirac partners $\chi_{B,W}$, since the gravitino mass $m_{3/2}$ has $R = -2$ in the sense of a spurion.

In the sequestered limit, we have $Z_1 \to 1$, $Z_2 \to 1$, and $Z_3 \to 0$, so that the eigenvalue and the eigenvector reduce to

$$m_\zeta \to 2m_{3/2}, \quad \bar{\Psi}_\zeta^T \to \begin{pmatrix} 0 & 0 & 0 & 0 & 0 \end{pmatrix}.$$  

(53)
as is expected. In other cases, the mass eigenvalue \( m_\xi \) is greatly suppressed, and, hence, the Higgsino component \( \tilde{\xi}^{(1)}_H \) is more suppressed than that of the gaugino Dirac partner, \( \tilde{\chi}^{(1)}_X \):

\[
\tilde{\xi}^{(1)}_H = \mathcal{O}\left( \frac{m_\xi}{m_{\text{weak}}} \right), \quad \tilde{\chi}^{(1)}_X = \mathcal{O}\left( \frac{m_{3/2}}{m_{\text{weak}}} \right),
\]

(54)

### 4.3. Implications for Higgs phenomenology and cosmology

Let us briefly discuss the implications for Higgs phenomenology and cosmology.

First, the invisible decay width of the 125 GeV Higgs into a pair of pseudo goldstinos is highly suppressed due to the approximate \( R \) symmetry. This is because fields with \( R \)-charge assignments other than \(+1\) contain a small goldstino component with suppression factor \( m_{2/3}/m_{\text{weak}} \). The Higgs decay into a pseudo goldstino and other neutralino is kinematically forbidden if the neutralino other than the pseudo goldstino is heavier than the lightest Higgs boson.

The second remark is concerned with the invisible cascade decay of the 125 GeV Higgs. If the scalar partner \( \phi \) of the pseudo goldstino is so light that the decay mode of the Higgs into a pair of \( \phi \) is kinematically allowed, the cascade decay into gravitinos and pseudo goldstinos \( h^0 \rightarrow \phi \phi \rightarrow \xi \tilde{G} \xi \tilde{G} \) dominates almost the entire Higgs decay channel [6], which contradicts the recent LHC discovery [40]. In our case, the scalar goldstino \( \phi \) receives a mass \( \lambda v^2 \) from the visible-sector SUSY breaking, in addition to the hidden-sector soft mass. Therefore, the invisible cascade decay of the Higgs is kinematically forbidden, if we assume \( \lambda > 0.36 \), or if the scalar pseudo goldstino has a soft mass larger than half of the Higgs mass.

Next we briefly discuss the cosmological constraints on the present model. Since we are supposing low-scale SUSY breaking in the visible sector, the lightest observable-sector supersymmetric particle (LOSP) mainly decays into a pseudo goldstino rather than a gravitino. This fact helps us avoid the constraint from BBN [11].

We should still worry about the overproduction problem, both for the gravitino and the pseudo goldstino. Let us first discuss the gravitino case. If gravitinos are thermally produced in the early Universe, the estimated abundance easily exceeds the limit \( \Omega h^2 < 0.1 \) for gravitinos heavier than \( \mathcal{O}(100 \text{ eV}) \). On the other hand, if gravitinos are not thermalized [41], overproduction excludes masses smaller than 10 keV, whereas gravitinos heavier than \( \mathcal{O}(10 \text{ keV}) \) can avoid overproduction if the reheating temperature is sufficiently low. Moreover, for masses smaller than \( \mathcal{O}(100 \text{ eV}) \), the allowed region is greatly reduced to \( m_{3/2} \lesssim 16 \text{ eV} \) due to the warm dark matter constraint [43,44]. Putting all this together, we must assume that the gravitino mass satisfies either \( m_{3/2} \lesssim 16 \text{ eV} \) or \( 10 \text{ keV} \lesssim m_{3/2} \), as was announced in Eq. (1).

In contrast, the pseudo goldstino couples fairly strongly to the MSSM particles and is expected to be thermalized. Recalling that its mass \( m_\xi \) is greatly suppressed compared to the gravitino mass \( m_{3/2} \) as in Eq. (38), the pseudo goldstino will be lighter than 16 eV for a gravitino mass \( m_{3/2} \gtrsim 10 \text{ keV} \) depending on the suppression factor \( Z_1 \), it may be possible that the pseudo goldstino avoids the overclosure problem even when the gravitino mass is as large as \( \mathcal{O}(10 \text{ MeV}) \). If so, the pseudo goldstino can make only a small contribution to the relic density.

In the present model, the lightest supersymmetric particle (LSP) is the pseudo goldstino whereas the next-to-LSP (NLSP) is the gravitino, which can decay into a photon and a pseudo goldstino. In

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\(^{12}\)If the reheating temperature is low, thermal leptogenesis is difficult to achieve. The possibility of electroweak baryogenesis in \( R \)-symmetric models was recently discussed in Ref. [42].
this situation, we should also be careful about the constraints from the gamma-ray line search from the Galactic Center region and the overproduction of the isotropic diffuse photon background; there is a stringent limit [45] on mono-energetic photons emitted from decaying particles. In our case, we expect that such gravitino decay is greatly suppressed since the pseudo goldstino mass eigenstate has very small gaugino and Higgsino components. If this is really the case, the gravitino is sufficiently long-lived and the present model can avoid the constraint from the diffuse gamma-ray line search. A preliminary consideration shows that only gravitinos heavier than \(O(100 \text{ MeV})\) would be excluded, even if the mixing between the photino and the pseudo goldstino is of order \(O(m_{3/2}/m_{\text{weak}})\).

5. Conclusion and discussion

We have studied the \(U_R(1)\)-symmetric extension of the SUSY standard model, which contains Dirac gauginos and an extended Higgs sector, assuming that the Dirac gaugino masses are induced from hidden-sector SUSY breaking while the extended Higgs sector incorporates the visible SUSY breaking. The \(U(1)_R\) symmetry in the visible sector is broken solely by the minimal coupling to supergravity, so that \(U(1)_R\)-breaking effects are proportional to the gravitino mass. After the EWSB, the \(R\)-charged Higgs fields develop nonzero VEVs due to the tadpole terms induced from the \(U(1)_R\) breaking. The VEVs are proportional to the gravitino mass and hence can be very small when the gravitino mass is much smaller than the weak scale.

We have then shown the generation of small neutrino masses through the \(U(1)_R\)-breaking effects. With our flipped \(U(1)_R\) charge assignment, the right-handed neutrinos couple to the \(R\)-partner Higgs field \(X_u\), which develops a very small VEV. Therefore, small neutrino masses can naturally be obtained even if the neutrino Yukawa couplings are of order unity, provided that the gravitino mass is as low as \(1\)–\(10 \text{ eV}\). We have also examined how our generation mechanism for the neutrino mass is constrained from the proton decay induced by the dimension-five operator, \(U^c D^c D^c N^c\), allowed by our \(U(1)_R\) charge assignment. Interestingly, the amplitude for the dominant decay mode involves neutrino Yukawa coupling, instead of the top Yukawa. Therefore, the constraint from proton stability is severer for a larger neutrino Yukawa coupling, which implies a small gravitino mass for a fixed neutrino mass. We have estimated the lifetime of protons and found that the gravitino should be heavier than \(1 \text{ keV}\) if we require the coefficient \(C_5\) of the dimension-five operator to be of order unity and the cutoff \(\Lambda_{\text{cutoff}}\) to be the Planck scale. Actually, a gravitino mass between \(10 \text{ eV} < m_{3/2} < 10 \text{ keV}\) is cosmologically excluded [43,44,46]; hence, our model for the neutrino mass generation evades the constraint from proton decay and is cosmologically safe if the gravitino is as heavy as \(10\)–\(100 \text{ keV}\). We have also suggested how the constraint from the proton stability can be relaxed if we adopt \(U(1)\) flavor symmetry for generating the hierarchical structure of Yukawa matrices.

Another significant effect of the \(U(1)_R\) breaking is the mass and the mixing of the pseudo goldstino, which is a Nambu–Goldstone fermion of the visible SUSY breaking and is massless in the \(U(1)_R\)-symmetric limit. We have analyzed the neutralino mass matrix that incorporates the Dirac mass parameters of gauginos and Higgsinos, and obtained to the lowest order in \(R\) breaking the smallest mass eigenvalue and the corresponding mixing angles to other neutralinos. According to the general argument, the mass of the pseudo goldstino in the sequestered limit is twice the gravitino mass. In our case, the pseudo goldstino has a coupling to the Higgs fields, and, consequently, its mass is proportional to but suppressed from the gravitino mass. In our concrete calculation, in which the pseudo goldstino gets a mass from the contact term in the Kähler potential, the suppression factor is very small if the cutoff scale \(\Lambda_0\) of the contact term is much larger than the weak scale. For the
mixings, we have found that the pseudo goldstinos have suppressed contaminations of gauginos and Higgsinos.

As we mentioned above, a gravitino as heavy as 100 keV is cosmologically safe and also satisfies the proton stability. In this case, the gravitino is the NLSP while the LSP is the pseudo goldstino, which can be as light as 10 eV, or even much lighter. It is then tempting to speculate that the gravitino could be part of the cold dark matter, and that the pseudo goldstino could play a role of dark radiation. To examine such possibilities, we should investigate the properties of the pseudo goldstino more extensively, for instance, by including the anomaly-mediated gaugino masses and other quantum corrections to the mass matrix. This will also be important for a quantitative study of the diffuse gamma-ray line.

Although we have focused in the present work on the constraints from the proton stability and also on some cosmological ones, there remain many issues to be discussed: the constraints from EW precision measurements and an implementation of the 125 GeV Higgs should be examined more intensively; the phenomenological impacts of large neutrino Yukawa couplings, as in other “neutrino-phobic” models, may deserve further study in our $U(1)_R$-symmetric model. The assumptions behind our present setup should also be elucidated: a partial list includes the origin of the hidden-sector SUSY breaking, UV completion of the visible SUSY breaking sector, and the embedding into a semi-simple gauge group. In particular, the smallness of the gravitino mass should be examined in a consistent framework, although we have treated it simply as a free parameter. In this respect, it may be possible that the smallness of the gravitino mass as well as the absence of certain operators can be explained in the (warped) extra-dimensional setup along the lines of Ref. [47].

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Appendix A. Baryon- and lepton-number-violating operators

In this appendix, we systematically study baryon- and/or lepton-number-violating operators allowed by the gauge and $U(1)_R$ symmetries of the present model.

We first categorize the fields with the same gauge quantum numbers as follows:

$$ S \in \{ X_0, N^c \}, \quad \Phi_u \in \{ H_u, X_u \}, \quad \Phi_d \in \{ H_d, X_d, L \}. \quad (A1) $$

The operators allowed by the gauge symmetries can be classified according to their mass dimensions as

$$ W_2 = S, \quad (A2) $$

$$ W_3 = SW_2 + \Phi_u \Phi_d, \quad (A3) $$

$$ W_4 = SW_3 + (Q \Phi_d) D^c + (\Phi_d \Phi_d) E^c + U^c D^c D^c, \quad (A4) $$

$$ W_5 = SW_4 + (QQ)(Q \Phi_d) + U^c U^c D^c D^c + (Q \Phi_d) U^c E^c + (\Phi_u \Phi_d) (\Phi_u \Phi_d) + (QQ) U^c D^c. \quad (A5) $$
where we suppress coefficients and flavor indices for simplicity. We also omit operators involving the adjoint chiral fields since they have no effect on proton decay.

It is straightforward to see that all of the operators violating either the baryon or lepton number in $W_2$, $W_3$, and $W_4$ are forbidden by the $R$ symmetry: this includes the usual trilinear terms, $U^c D^c D^c$, $Q D^c L$, $L E^c L$, and the bilinear terms $H_u L$. We also note that the mass terms $X_0 N^c$ as well as $N^c N^c$ are forbidden by the $U(1)_R$. Therefore we can define the conserved baryon and lepton numbers at the renormalizable level.

There are three types of dimension-five operators that violate both the baryon and lepton numbers: $(Q Q)(Q L)$, $U^c U^c D^c E^c$, and $N^c U^c D^c D^c$. Among them, the first two are forbidden by the $U(1)_R$, whereas the last one is $U(1)_R$ symmetric in our flipped assignment. As a result, we are left with a unique operator $N^c U^c D^c D^c$, which is the baryon- and lepton-number-violating operator in the present model.

As we mentioned in Sect. 3.1, $(L H_u)(L H_u)$ operators are allowed, but they make a negligible contribution to neutrino masses if the cutoff $\Lambda_{\text{cutoff}}$ is sufficiently large. We also note that there are dimension-five operators, $X_0^2 N^c N^c$ and $(X_H X_d) N^c N^c$, which are $R$-invariant and lepton-number violating. After the $R$-Higgses get the VEVs, these operators lead to Majorana masses for right-handed neutrinos, of order $m_{3/2}^2/\Lambda$, which can be safely neglected (unless $m_{3/2} \sim 100$ TeV).

**Appendix B. $U(1)$ flavor symmetry and proton stability**

As we see in Sect. 3, the model is severely constrained from proton decay if the gravitino is very light, $m_{3/2} < 16$ eV, especially for large tan $\beta$. In this appendix, we consider a model with Froggatt–Nielson $U(1)$ flavor symmetry to relax the constraint.

Let $\Theta$ be a Froggatt–Nielson (FN) field with $U(1)_F$ charge $-1$, which is assumed to develop a nonzero VEV $(\Theta)$. We also denote the $U(1)_F$ charge of the fields $\Phi_i$ by the lower-case letter $\phi_i$. The Yukawa terms in Eq. (18) should be multiplied by a suitable power of the FN field $\Theta$; for instance, the neutrino Yukawa term is generated from the $U(1)_F$-invariant operator

$$W = f^{ij}_N \left( \frac{\Theta}{\Lambda_{\text{cutoff}}} \right)^{n_i^c + \ell_j + X_u} N_i^c L_j X_u,$$

where $f^{ij}_N$ is a coefficient of order 1 and $\Lambda_{\text{cutoff}}$ is the cutoff. After we substitute the FN field by its VEV, we have $y^{ij}_N = f^{ij}_N e^{\nu_i^c + \ell_j + X_u}$. We will identify $\varepsilon \equiv (\Theta)/\Lambda_{\text{cutoff}}$ with the Wolfenstein parameter $\varepsilon = 0.22$. Similarly, we have

$$y^{ij}_U = f^{ij}_U e^{
u_i^c + w_j^c + h_u}, \quad y^{ij}_D = f^{ij}_D e^{\nu_i^c + d_j^c + h_d}, \quad y^{ij}_E = f^{ij}_E e^{\nu_i^c + \ell_j + h_d}. \quad (B1)$$

The $U(1)_F$ charge assignment can be restricted by the following requirements. First, from the intergenerational mixings of quarks, $V_{us} \sim \varepsilon$, $V_{cb} \sim \varepsilon^2$, and $V_{ub} \sim \varepsilon^3$, we obtain the following relations:

$$q_1 = q_3 + 3, \quad q_2 = q_3 + 2. \quad (B2)$$

Next, the quark mass hierarchy can be parameterized as $m_u : m_c : m_t \sim \varepsilon^8 : \varepsilon^4 : 1$ and $m_d : m_s : m_b \sim \varepsilon^4 : \varepsilon^2 : 1$, while the mass hierarchy of charged leptons is $m_e : m_\mu : m_\tau \sim \varepsilon^5 : \varepsilon^2 : 1$. These mass hierarchies can be reproduced, for instance, if we take

$$u_1^c = u_5^c + 5, \quad d_1^c = d_3^c + 1, \quad \ell_1 + e_1^c = \ell_3 + e_5^c + 5,$$

$$u_2^c = u_5^c + 2, \quad d_2^c = d_3^c, \quad \ell_2 + e_2^c = \ell_3 + e_5^c + 2. \quad (B3)$$
The FN charges of the third generation of quarks and leptons can be constrained, if we require top Yukawa coupling of order unity, \( y_t \sim 1 \), and also bottom-tau unification, \( m_b \sim m_{\tau} \), which gives

\[
q_3 + u_3^c + h_u = 0, \quad q_3 + d_3^c = e_3^c + \bar{\ell}_3.
\]  

(B4)

In addition, we require for definiteness that the supersymmetric Higgs mass terms have null charge,

\[
x_d + h_u = 0, \quad x_u + h_d = 0.
\]  

(B5)

Note that the FN charges of \( H_u \) and \( H_d \) can be chosen independently since there is no \( \mu H_u H_d \) term in the present \( U(1)_R \)-symmetric model.

Now, we discuss the suppression of the dimension-five operator (22) by the FN mechanism. The coefficient \( C_5 = C_5^{1123} \) of the operator \( U_1^c D_1^c D_2^c N_3^c \) receives a suppression of \( C_5 \sim \varepsilon^{u_1^c + d_1^c + d_3^c + e_3^c} \).

The FN charge \( n_c^5 \) of the right-handed neutrino can be eliminated if we use the relation (26)

\[
\nu_{ij} = f_{ij}^e n_c^1 + l_j + x_u = \frac{m_{\nu_j}^2}{m_{3/2}} \left( \frac{\mu_2^2 + m_{X_u}^2}{\mu_2 \cos \beta} \right).
\]  

(B6)

Combining Eqs. (B2)–(B5) with Eq. (B6), we obtain

\[
C_5 = C_5^{(0)} e^{u_1^c + d_1^c + d_3^c + e_3^c} = C_5^{(0)} \varepsilon^6 + Q \left( \frac{m_{\nu_j}^2}{m_{3/2}} \right) \left( \frac{\mu_2^2 + m_{X_u}^2}{\mu_2 \cos \beta} \right),
\]  

(B7)

where

\[
Q \equiv h_d - h_u + d_3^c + e_3^c - 2q_3.
\]  

(B8)

We see that the coefficient \( C_5 \) is suppressed for several reasons: the suppression factor \( \varepsilon^6 \) comes from the FN charges of the first and second generations of quarks involved in the proton decay. Another factor \( \varepsilon^Q \) expresses the dependence on the FN charges of the Higgses and the third generation of quarks and leptons. For instance, if we assign a negative FN charge to the MSSM Higgs \( H_u \), the up-type quarks have positive FN charges, which suppress the dimension-five operator. Notice that this can be done without suppressing the supersymmetric mass terms since we can assign the opposite FN charge to the \( R \)-partner \( X_d \).

It should be noticed that there is yet another suppression in Eq. (B7): the coefficient \( C_5 \) is more suppressed for larger gravitino masses. This dependence can be understood as follows. Assigning a positive FN charge \( n_c^e \) to 0 to the right-handed neutrinos suppresses not only the coefficient \( C_5 \) of the dimension-five operator but also the neutrino Yukawa coupling \( y_\nu \); \( y_\nu \sim \varepsilon^{n_c^e} \) and \( C_5 \sim \varepsilon^{n_c^e} \), so that the proton lifetime scales as \( \tau_p \sim e^{-4 n_c^e} \). For a fixed value of the neutrino mass, suppression of the neutrino Yukawa coupling is translated into dependence on \( m_{3/2} \) in Eq. (B7).

Putting all this together, a similar analysis as in Sect. 3.2 gives an estimate of the proton lifetime as

\[
\tau_p = \frac{3.7 \times 10^{34}}{\varepsilon^2 Q} \text{yr} \left( \frac{\Lambda_{\text{cutoff}}}{2.4 \times 10^{18} \text{ GeV}} \right)^2 \left( \frac{m_{3/2}}{1 \text{ eV}} \right)^4 \left( \frac{\cos^2 \beta}{0.1} \right)^3,
\]  

(B9)

where we have set the coefficient to be \( C_5^{(0)} = 1 \). Figure B1 shows how the constraint from the proton decay can be satisfied for a moderate choice of the FN charge \( Q \). We see that the proton decay constraint can be satisfied even for \( m_{3/2} = 1 \text{ eV} \) if we take \( Q = 0 \) and \( \tan \beta = 3 \). In this way, the present model combined with the \( U(1)_F \) flavor symmetry evades the proton decay constraint in a wider region of parameter space.
The situation is quite different from that in the MSSM case [39], in which a negative FN charge of the Higgses would imply an unacceptably small $\mu$ term. In fact, the analysis in Ref. [39] shows that the FN charge of the $\mu$ term, i.e., the sum of the FN charges of $H_u$ and $H_d$, is tightly constrained since both the $LLLL$ and $RRRR$ operators (31) should be suppressed simultaneously. In our case, we only have the $RRRR$-type operator since the $LLLL$ operator is forbidden by the $U(1)_R$ symmetry.

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