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R-symmetric flipped SU(5)

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ABSTRACT: We construct a supersymmetric flipped SU(5) grand unified model that possesses an *R* symmetry. This *R* symmetry forbids dangerous non-renormalizable operators suppressed by a cut-off scale up to sufficiently large mass dimensions so that the SU(5)-breaking Higgs field develops a vacuum expectation value of the order of the unification scale along the *F*- and *D*-flat directions, with the help of the supersymmetry-breaking effect. The mass terms of the Higgs fields are also forbidden by the *R* symmetry, with which the doublet-triplet splitting problem is solved with the missing partner mechanism. The masses of right-handed neutrinos are generated by non-renormalizable operators, which then yield a light neutrino mass spectrum and mixing through the seesaw mechanism that are consistent with neutrino oscillation data. This model predicts one of the color-triplet Higgs multiplets to lie at an intermediate scale, and its mass is found to be constrained by proton decay experiments to be $\gtrsim 5 \times 10^{11}$ GeV. If it is $\lesssim 10^{12}$ GeV, future proton decay experiments at Hyper-Kamiokande can test our model in the $p \rightarrow \pi^0 \mu^+$ and $p \rightarrow K^0 \mu^+$ decay modes, in contrast to ordinary grand unified models where $p \rightarrow \pi^0 e^+$ or $p \rightarrow K^+ \bar{\nu}$ is the dominant decay mode. This characteristic prediction for the proton decay branches enables us to distinguish our model from other scenarios.

KEYWORDS: Supersymmetry Phenomenology

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1 Introduction

Quarks and leptons in the Standard Model (SM) may be unified at high energies. Since 1970's, there have been many efforts to construct a concrete model that accommodates the unification of quarks and leptons, as well as of the strong and electroweak interactions [1–6]. Most of these models — dubbed as grand unified theories (GUTs) — predict that the unification is achieved at a very high energy [7], and the large difference between the electroweak and unification scales brings about the hierarchy problem [8–12]. The electroweak scale may be stabilized against this large hierarchy if supersymmetry (SUSY) appears above the electroweak scale, which highly motivates SUSY GUTs [13, 14]. In particular, the minimal SUSY SU(5) GUT [13, 14] has been regarded as a representative, promising model of GUTs as the matter content in the minimal SUSY SM (MSSM) nicely fits SU(5) representations and, among other things, gauge coupling unification is found to be achieved with great accuracy [15–21].

A drawback of the minimal SU(5) is that a high level of fine-tuning is required to keep the MSSM Higgs fields from acquiring a GUT-scale mass while making their SU(5) partner fields massive. This problem, called the doublet-triplet splitting problem, is a generic problem in GUTs, and many solutions to this problem have been proposed so far. Among them, the missing partner mechanism [22, 23] is most frequently discussed in SU(5) GUTs, where a **75** and a pair of **50** and $\overline{\mathbf{50}}$ representations of SU(5) are introduced in addition to

the $\mathbf{10}$ and $\bar{\mathbf{5}}$ representations for the MSSM matter fields and $\mathbf{5}$ and $\bar{\mathbf{5}}$ for the MSSM Higgs fields. The $SU(5)$ gauge symmetry is spontaneously broken by a vacuum expectation value (VEV) of the $\mathbf{75}$ field and it also gives vector-like masses to the color-triplet components in the $\mathbf{5}$ and $\bar{\mathbf{5}}$ Higgs fields together with the corresponding components in the $\mathbf{50}$ and $\bar{\mathbf{50}}$ fields, whereas the doublet components in the $\mathbf{5}$ and $\bar{\mathbf{5}}$ Higgs fields, corresponding to the MSSM Higgs fields, do not acquire masses from the VEV of the $\mathbf{75}$. In order for this mechanism to work, the absence of the bilinear term for the $\mathbf{5}$ and $\bar{\mathbf{5}}$ Higgs fields is crucial, which may be attributed to an additional symmetry [24–29]. In spite of the conceptual success of the missing partner mechanism, it is known that the missing partner model with the minimal matter content is incompatible with the perturbative gauge coupling unification due to the existence of large representations [30] and thus a more elaborate model building is required to realize this mechanism in $SU(5)$.

Another economical way to realize this missing partner mechanism is found in the flipped $SU(5) \times U(1)$ model [31–34], where the right-handed charged leptons and neutrinos, and the right-handed up- and down-type quarks, are flipped with respect to the standard $SU(5)$ assignment. In this model, the $SU(5) \times U(1)$ gauge symmetry is broken by the VEVs of $\mathbf{10}$ and $\bar{\mathbf{10}}$ Higgs fields, and they give masses only to the color-triplet components in the $\mathbf{5}$ and $\bar{\mathbf{5}}$ Higgs fields [33]. As a result, if there is no bilinear term of the $\mathbf{5}$ and $\bar{\mathbf{5}}$ Higgs fields, then the MSSM Higgs fields remain massless after the $SU(5) \times U(1)$ symmetry is spontaneously broken. Another feature of this model is that the $\mathbf{10}$ and $\bar{\mathbf{10}}$ Higgs fields develop VEVs along the F - and D -flat directions with the help of the effect of soft masses as well as non-renormalizable operators. Such non-renormalizable operators must be suppressed up to sufficiently large mass dimensions in order for these VEVs to be of the order of the unification scale. The suppression of these operators, as well as the absence of the bilinear term of the $\mathbf{5}$ and $\bar{\mathbf{5}}$ Higgs fields, are just assumed in the standard flipped $SU(5)$ model [33] — it is, therefore, desirable to account for these key assumptions by means of certain symmetries.

In this work, we construct a flipped $SU(5)$ model that can resolve this issue. To that end, we require the model to have a global $U(1)_R$ symmetry which forbids all of the unwanted terms mentioned above. This R symmetry is supposed to be explicitly broken only by the constant term in the superpotential so that we can obtain a viable SUSY mass spectrum as well as the vanishing cosmological constant. We then show that the $U(1)_R$ and $SU(5) \times U(1)$ symmetries are broken at the scales of $\simeq 3 \times 10^{17}$ GeV and $\simeq 10^{16}$ GeV, respectively, and that the μ -term of the MSSM Higgs doublets is naturally suppressed thanks to the missing partner mechanism.

As it turns out, this model predicts one of the color-triplet Higgs multiplets to lie at an intermediate scale. We evaluate the current limit on the mass of this light color-triplet Higgs from proton decay experiments and discuss its implications for the future experiments at Hyper-Kamiokande. It is found that in the presence of a color-triplet Higgs with a mass $\lesssim 10^{12}$ GeV, proton decay can be discovered at Hyper-Kamiokande in the $p \rightarrow \pi^0 \mu^+$ and $p \rightarrow K^0 \mu^+$ decay modes. Note that in ordinary GUT models either $p \rightarrow \pi^0 e^+$ or $p \rightarrow K^+ \bar{\nu}$ is the dominant decay mode. We may therefore distinguish our model from others with proton decay experiments, which highly motivates to search for

the $p \rightarrow \pi^0 \mu^+$ and $p \rightarrow K^0 \mu^+$ modes in future proton decay experiments, in addition to the standard channels $p \rightarrow \pi^0 e^+$ and $p \rightarrow K^+ \bar{\nu}$.

This paper is organized as follows. In the subsequent section, we describe our model and discuss its symmetry breaking structure and mass spectrum. In section 3, we study the flavor structure of this model, and show that the observed pattern of the SM quark and lepton masses and mixings can be reproduced. We then compute the proton decay rates in section 4 and discuss the testability of our model. Finally, section 5 is devoted to conclusion and discussion.

2 R -symmetric flipped SU(5)

2.1 Model

The model discussed in this paper is based on an $SU(5) \times U(1)$ gauge theory [31–34], where the three generations of MSSM matter fields, as well as three right-handed neutrino chiral superfields, are embedded into $\mathbf{10}(1)$, $\bar{\mathbf{5}}(-3)$ and $\mathbf{1}(5)$ representations, with the numbers in the parentheses indicating the $U(1)$ charges in units of $1/\sqrt{40}$. We denote these representations by F_i , $\bar{f}_i$, and ℓ_i^c , respectively, with $i = 1, 2, 3$ the generation index. In addition to these matter fields, this model contains a pair of $\mathbf{10}(1)$ and $\bar{\mathbf{10}}(-1)$ Higgs fields, H and $\bar{H}$, a pair of $\mathbf{5}(-2)$ and $\bar{\mathbf{5}}(2)$ Higgs fields, h and $\bar{h}$, and a singlet field, S . The H and $\bar{H}$ fields break the $SU(5) \times U(1)$ gauge group down to the SM gauge group once these fields develop VEVs. As in the MSSM, the $SU(2)_L \times U(1)_Y$ gauge symmetry is broken by the VEVs of the doublet Higgs fields H_d and H_u , which reside in h and $\bar{h}$, respectively. In table 1, we summarize the field content and the charge assignments of the fields. With these charge assignments, the $U(1)_Y$ hypercharge Y is given by the following linear combination of the $SU(5)$ generator $T_{24} = \text{diag}(2, 2, 2, -3, -3)/\sqrt{60}$ and the $U(1)$ charge Q_X :¹

$$Y = \frac{1}{\sqrt{15}} T_{24} + \sqrt{\frac{8}{5}} Q_X. \quad (2.1)$$

The second column in table 1 summarizes the component fields in each representation; for instance, F_i and $\bar{f}_i$ contain the MSSM matter fields as²

$$F_i = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & d_{i3}^c & -d_{i2}^c & u_i^1 & d_i^1 \\ -d_{i3}^c & 0 & d_{i1}^c & u_i^2 & d_i^2 \\ d_{i2}^c & -d_{i1}^c & 0 & u_i^3 & d_i^3 \\ -u_i^1 & -u_i^2 & -u_i^3 & 0 & \nu_i^c \\ -d_i^1 & -d_i^2 & -d_i^3 & -\nu_i^c & 0 \end{pmatrix}, \quad \bar{f}_i = \begin{pmatrix} u_{i1}^c \\ u_{i2}^c \\ u_{i3}^c \\ e_i \\ -\nu_i \end{pmatrix}. \quad (2.2)$$

Note that the right-handed neutrino fields ν_i^c are necessary ingredients in flipped SU(5) GUT models to form complete $\mathbf{10}$ representations, in contrast to the standard SU(5) GUT [3], where right-handed neutrino fields are singlet under the SU(5) gauge group.

¹We adopt the normalization of hypercharge such that left-handed neutrinos have hypercharge $-1/2$. We also use the SU(5) normalization of hypercharge, $Q_1 \equiv \sqrt{3/5} Y$. The gauge coupling constants corresponding to Y and Q_1 are denoted by g' and g_1 , respectively.

²In these equations, we neglect the mixing among different generations just for simplicity. More complete expressions for the embedding of the component fields are given in section 3.

Fields	Components	SU(5)	U(1)	U(1) _R
F_i	d_i^c, Q_i, ν_i^c	10	+1	17/36
$\bar{f}_i$	u_i^c, L_i	$\bar{\mathbf{5}}$	−3	17/36
ℓ_i^c	e_i^c	1	+5	17/36
H	d_H^c, Q_H, ν_H^c	10	+1	1/36
$\bar{H}$	$d_{\bar{H}}^c, Q_{\bar{H}}, \nu_{\bar{H}}^c$	$\bar{\mathbf{10}}$	−1	17/36
h	D, H_d	5	−2	19/18
$\bar{h}$	$\bar{D}, H_u$	$\bar{\mathbf{5}}$	+2	19/18
S	S	1	0	1/9

Table 1. The field content and the charge assignments in our model. The U(1) charges are given in units of $1/\sqrt{40}$.

In addition, we assume that this model respects a global U(1)_R symmetry in the global SUSY limit. The U(1)_R charge assignment is also shown in table 1, where we normalize the U(1)_R charge such that the superpotential has the U(1)_R charge +2. As we see below, this R symmetry plays an important role in suppressing unwanted terms. We can instead consider a discrete R symmetry with the same matter content to suppress unwanted terms; we show an example for this case in appendix A. In either case, the R symmetry is broken by the VEV of the singlet field S .

We note that this model contains only one singlet field, S , contrary to the standard flipped SU(5) model [33, 35–38], where several singlet fields are introduced to generate neutrino masses via the double seesaw mechanism [39–43]; namely, our model has a simpler matter content than the ordinary one. In the present case, the masses of the right-handed neutrinos are generated via non-renormalizable operators [44], as we see below.

The superpotential terms allowed by these symmetries are

$$W = W_{\text{Yukawa}} + W_{\text{DT}} + W_{\text{neutrino}} + W_{HS} \dots, \quad (2.3)$$

with

$$W_{\text{Yukawa}} = \frac{1}{4} \lambda_1^{ij} \epsilon_{\alpha\beta\gamma\delta\epsilon} F_i^{\alpha\beta} F_j^{\gamma\delta} h^\epsilon + \sqrt{2} \lambda_2^{ij} F_i^{\alpha\beta} \bar{f}_{j\alpha} \bar{h}_\beta + \lambda_3^{ij} \bar{f}_{i\alpha} \ell_j^c h^\alpha, \quad (2.4)$$

$$W_{\text{DT}} = \frac{\lambda_4}{4\Lambda_{\text{DT}}^8} \epsilon_{\alpha\beta\gamma\delta\epsilon} S^8 H^{\alpha\beta} H^{\gamma\delta} h^\epsilon + \frac{1}{4} \lambda_5 \epsilon^{\alpha\beta\gamma\delta\epsilon} \bar{H}_{\alpha\beta} \bar{H}_{\gamma\delta} \bar{h}_\epsilon, \quad (2.5)$$

$$W_{\text{neutrino}} = \frac{c_{ij}}{2\Lambda_N^2} S (F_i^{\alpha\beta} \bar{H}_{\alpha\beta}) (F_j^{\gamma\delta} \bar{H}_{\gamma\delta}), \quad (2.6)$$

$$W_{HS} = \frac{\lambda_H}{4\Lambda_{HS}^5} (H^{\alpha\beta} \bar{H}_{\alpha\beta})^4 + \frac{\lambda_{HS}}{18\Lambda_{HS}^{10}} (H^{\alpha\beta} \bar{H}_{\alpha\beta})^2 S^9 + \frac{\lambda_S}{18\Lambda_{HS}^{15}} S^{18}, \quad (2.7)$$

where the Greek superscripts and subscripts denote the SU(5) indices, $\epsilon_{\alpha\beta\gamma\delta\epsilon}$ is the totally antisymmetric tensor, and the dots in eq. (2.3) indicate higher-dimensional operators that are irrelevant to the following discussions.³ Λ_{DT} , Λ_N , and Λ_{HS} are parameters with a mass

³Exceptions are those who cause R -parity violation and dimension-five proton decay operators at low energies, which we discuss in section 2.4 and in section 4, respectively.

dimension one that correspond to a cut-off scale of this model. Considering that the non-renormalizable operators in eqs. (2.5)–(2.7) depend on a large power of the cut-off scale and that different types of fields and interactions may be responsible for generating these operators, we have introduced different parameters for the cut-off scales of these operators, though we expect all of them to be around the reduced Planck mass scale, M_P . The terms in eq. (2.4) contain the MSSM Yukawa coupling terms. After the $SU(5) \times U(1)$ and $U(1)_R$ symmetries are broken, W_{DT} gives masses to the color-triplet Higgs fields and W_{neutrino} leads to a Majorana mass matrix for right-handed neutrinos, as we see in section 2.2. W_{HS} in eq. (2.7) includes all of the terms that are composed of the H , $\bar{H}$, and S fields and allowed by the symmetries of the theory; as we see, the $U(1)_R$ symmetry highly restricts possible forms of such operators. We take λ_H and λ_S to be real and positive without loss of generality. Notice that the bilinear terms $H\bar{H}$ and $h\bar{h}$ are both forbidden by the $U(1)_R$ symmetry.⁴ In the standard flipped $SU(5)$ model, the $H\bar{H}$ term is forbidden by a $\mathbb{Z}_2$ symmetry, while the $h\bar{h}$ term is generically allowed. As we see in section 2.2, the absence of the $h\bar{h}$ term is crucial to solve the doublet-triplet splitting problem.

In the following analysis, we assume that the coefficients of the operators in eqs. (2.5)–(2.7) are $\mathcal{O}(1)$; i.e., the hierarchy in the contribution of these operators is completely controlled by the symmetries of the model. For the terms in eq. (2.4), on the other hand, we allow the coefficients to be much smaller than $\mathcal{O}(1)$ to reproduce the SM Yukawa couplings. Although the structure of the SM Yukawa couplings may also be explained by using the $U(1)_R$ symmetry with a generation-dependent charge assignment — in a similar manner to the Froggatt-Nielsen model [45] — we do not pursue this possibility in the present paper and defer it to future work [46].

2.2 Symmetry breaking and mass spectrum

Next, we study the vacuum structure of this model. In the following analysis, we assume the canonical form of the Kähler potential, $K = |F_i^{\alpha\beta}|^2 + \dots$, just for simplicity. In this case, the scalar potential of the theory is computed as

$$V = V_F + V_D + V_{\text{soft}}, \quad (2.8)$$

with

$$V_F = \sum_i \left| \frac{\partial W}{\partial \phi_i} \right|^2, \quad (2.9)$$

$$V_D = \frac{g_5^2}{2} \left[\sum_i \phi^{*i} T_i^A \phi_i \right]^2 + \frac{g_X^2}{2} \left[\sum_i Q_{X_i} |\phi_i|^2 \right]^2, \quad (2.10)$$

where ϕ_i denote the scalar components of the chiral superfields in this model, g_5 and g_X are the gauge coupling constants of the $SU(5)$ and $U(1)$ gauge interactions, respectively, T_i^A and Q_{X_i} are the representation of the $SU(5)$ generators and $U(1)_X$ charge of the field ϕ_i , respectively, and V_{soft} denotes the contribution of soft SUSY-breaking terms to the scalar potential.

⁴The terms $F_i \bar{H}$ are also prohibited by the $U(1)_R$ symmetry.

We search for a minimum of the potential at which ν_H^c , $\nu_{\bar{H}}^c$, and S develop VEVs;⁵ the SM gauge group is unbroken in this case. In fact, it turns out that there is a F - and D -flat direction for these fields in the absence of SUSY-breaking effect and non-renormalizable operators. To see this, let us set all of the fields other than these three fields to be zero and $V_{\text{soft}} = 0$; in this case, at renormalizable level, $V_F = 0$ and

$$V_D = \left(\frac{3g_5^2}{10} + \frac{g_X^2}{80} \right) \left[|\nu_H^c|^2 - |\nu_{\bar{H}}^c|^2 \right]^2. \quad (2.11)$$

Therefore, the potential vanishes for $|\nu_H^c| = |\nu_{\bar{H}}^c|$ and arbitrary values of S .

We now include the effect of non-renormalizable operators and soft SUSY breaking masses, with which the potential terms for the above fields are

$$\begin{aligned} V_F = & \left| \frac{\lambda_H}{\Lambda_{HS}^5} (\nu_H^c \nu_{\bar{H}}^c)^2 + \frac{\lambda_{HS}}{9\Lambda_{HS}^{10}} S^9 \right|^2 |\nu_H^c \nu_{\bar{H}}^c|^2 \left(|\nu_H^c|^2 + |\nu_{\bar{H}}^c|^2 \right) \\ & + \left| \frac{\lambda_{HS}}{2\Lambda_{HS}^{10}} (\nu_H^c \nu_{\bar{H}}^c)^2 + \frac{\lambda_S}{\Lambda_{HS}^{15}} S^9 \right|^2 |S|^{16}, \end{aligned} \quad (2.12)$$

$$V_{\text{soft}} = -m_H^2 |\nu_H^c|^2 - m_{\bar{H}}^2 |\nu_{\bar{H}}^c|^2 - m_S^2 |S|^2, \quad (2.13)$$

where we assume m_H^2 , $m_{\bar{H}}^2$, and m_S^2 are all positive. Let us search for a minimum of the potential V in eq. (2.8) in the flat direction of the D -term potential in eq. (2.11); i.e., $|\nu_H^c| = |\nu_{\bar{H}}^c|$. To that end, we first note that by using the $U(1)$ gauge and $U(1)_R$ transformations, we can always take the VEV of S to be real and positive, and the VEVs of ν_H^c and $\nu_{\bar{H}}^c$ to satisfy $\langle \nu_H^c \rangle = \langle \nu_{\bar{H}}^c \rangle = \Phi$. The field Φ is sometimes called flaton [35–38]. Now recall that we have taken λ_H and λ_S to be real and positive. To estimate the size of the VEVs of Φ and S with a simple analysis, we also take λ_{HS} to be real and positive. In this case, we can express the potential in terms of $\Phi \equiv |\Phi|e^{i\theta}$ and S as

$$\begin{aligned} V = & 2|\Phi|^6 \left[\frac{\lambda_H^2}{\Lambda_{HS}^{10}} |\Phi|^8 + \frac{\lambda_{HS}^2}{81\Lambda_{HS}^{20}} S^{18} + \frac{2\lambda_H \lambda_{HS}}{9\Lambda_{HS}^{15}} S^9 |\Phi|^4 \cos(4\theta) \right] \\ & + S^{16} \left[\frac{\lambda_{HS}^2}{4\Lambda_{HS}^{20}} |\Phi|^8 + \frac{\lambda_S^2}{\Lambda_{HS}^{30}} S^{18} + \frac{2\lambda_{HS} \lambda_S}{2\Lambda_{HS}^{25}} S^9 |\Phi|^4 \cos(4\theta) \right] \\ & - m_\Phi^2 |\Phi|^2 - m_S^2 S^2, \end{aligned} \quad (2.14)$$

with $m_\Phi^2 \equiv m_H^2 + m_{\bar{H}}^2$. As we see, the potential minimum is obtained for $\cos(4\theta) = -1$, meaning that at the minimum of the potential, Φ^4 is real and negative. We then solve the minimization conditions of the potential for a positive value of S and a negative value of Φ^4 to determine their VEVs.

In figure 1, we show the absolute values of the VEVs $\langle S \rangle$ and $\langle \Phi \rangle$ at the minimum of the potential (2.14) as functions of $m_{\text{SUSY}} = m_S = m_\Phi$ in the blue solid and green dashed

⁵We assume that the soft mass squared terms of the MSSM sector are all positive except those in the Higgs sector, in order to avoid unwanted color and electromagnetic vacua; this situation may naturally be realized through renormalization group effect, since the soft masses of SM singlet fields decrease at low energies due to the absence of the contribution by gauge interactions [33].

Figure 1. The absolute values of the VEVs $\langle S \rangle$ and $\langle \Phi \rangle$ as functions of $m_{\text{SUSY}} = m_S = m_\Phi$ in the blue solid and green dashed lines, respectively, where we take $\lambda_H = \lambda_{HS} = \lambda_S = 1$ and $\Lambda_{HS} = M_P$. Each band shows the range of the predicted values when we vary the parameter Λ_{HS} from $M_P/2$ to $2M_P$.

lines, respectively, where we take $\lambda_H = \lambda_{HS} = \lambda_S = 1$ and $\Lambda_{HS} = M_P$. Each band shows the range of the predicted values when we vary the parameter Λ_{HS} from $M_P/2$ to $2M_P$. We see that $|\langle \Phi \rangle| \simeq 10^{16}$ GeV and $|\langle S \rangle| \simeq \text{a few} \times 10^{17}$ GeV in the range of m_{SUSY} shown in the plot. The dependence of these VEVs on the soft masses is mild since the terms in V_F are composed of a large power of the fields Φ and S .

Notice that if there were operators of H , $\bar{H}$, and S that have mass dimensions lower than those in eq. (2.7), such as $(H\bar{H})^2$, the VEVs $|\langle \Phi \rangle|$ and $|\langle S \rangle|$ would be much smaller than those in figure 1 and thus phenomenologically disfavored. In the ordinary flipped SU(5) models, such dangerous operators are just assumed to be absent, whilst in our model their absence is attributed to the $U(1)_R$ symmetry.

The VEV $\langle \Phi \rangle = \langle \nu_H^c \rangle = \langle \nu_{\bar{H}}^c \rangle$ breaks the $SU(5) \times U(1)$ gauge group into the SM gauge group. The number of broken gauge symmetries is $25 - (8 + 3 + 1) = 13$, and thus among the 10×2 components of H and $\bar{H}$, 13 components are absorbed by the gauge bosons, with 12 gauge bosons acquiring a mass M_{V_5} and the remaining one a mass M_{V_1} . These masses are computed as

$$M_{V_5} = g_5 |\langle \Phi \rangle|, \quad M_{V_1} = \sqrt{\frac{5}{2}} \left(\frac{24}{25} g_5^2 + \frac{1}{25} g_X^2 \right)^{\frac{1}{2}} |\langle \Phi \rangle|. \quad (2.15)$$

The remaining 7 components of H and $\bar{H}$ appear as physical states. One of them corresponds to the flaton field Φ , which acquires an $\mathcal{O}(m_{\text{SUSY}})$ mass at the minimum of the potential. The other six components are d_H^c and $d_{\bar{H}}^c$ in H and $\bar{H}$, respectively. After

ν_H^c and $\nu_{\bar{H}}^c$ develop a VEV, d_H^c and $d_{\bar{H}}^c$ form vector-like mass terms with D and $\bar{D}$ in h and $\bar{h}$ via the first and second terms in eq. (2.5), with the masses

$$M_{H_C} = \lambda_4 |\langle \Phi \rangle| \left(\frac{\langle S \rangle}{\Lambda_{\text{DT}}} \right)^8 \simeq \lambda_4 \times 7 \times 10^{11} \times \left(\frac{|\langle \Phi \rangle|}{10^{16} \text{ GeV}} \right) \left(\frac{\langle S \rangle}{3 \times 10^{17} \text{ GeV}} \right)^8 \left(\frac{\Lambda_{\text{DT}}}{10^{18} \text{ GeV}} \right)^{-8} \text{ GeV}, \quad (2.16)$$

$$M_{\bar{H}_C} = \lambda_5 |\langle \Phi \rangle|, \quad (2.17)$$

respectively. As we see, $M_{H_C} \ll M_{\bar{H}_C}$ in our model; this can be contrasted with the case in the standard flipped SU(5), where $M_{H_C} \simeq M_{\bar{H}_C}$. The implications of the presence of a light color-triplet Higgs multiplet will be discussed below.

Notice that even after the $\text{SU}(5) \times \text{U}(1)$ gauge group is broken, the MSSM doublets in h and $\bar{h}$ do not acquire GUT-scale masses and only the color-triplet components become massive. This shows that the doublet-triplet splitting problem is solved in this setup by the missing partner mechanism [22, 23, 33]. We also note that for this mechanism to work, the absence of large μ -term for h and $\bar{h}$ is crucial. In the current setup, the $\text{U}(1)_R$ symmetry prohibits not only the bilinear operator $h\bar{h}$, but also all of the operators of the form $S^m (H\bar{H})^n h\bar{h}$, and thus there is no danger for the generation of a large μ -term. This is another advantage in our model.

The VEVs of the Φ and S fields also break the $\text{U}(1)_R$ symmetry. Since this is a global symmetry, a Nambu-Goldstone boson, called R -axion [47], appears at low energies, which is a linear combination of the pseudo-scalar components of the $\text{U}(1)_R$ -charged fields that develop VEVs. If this $\text{U}(1)_R$ symmetry is an exact symmetry, this R -axion is massless. In many realistic SUSY-breaking scenarios, however, the $\text{U}(1)_R$ symmetry is explicitly violated by a constant term in the superpotential, which is used to suppress the cosmological constant. We also assume an $\text{U}(1)_R$ -breaking constant term in the superpotential, W_{const} , in our model.⁶ Let us write the superpotential as $W = W_0 + W_{\text{const}}$, where W_0 denotes the superpotential that describes both the visible and hidden sectors. In the presence of W_{const} , the R -axion acquires a mass through the supergravity effect. At the leading order in the M_P^{-1} expansion, it is computed as [48]

$$m_a^2 \simeq \frac{8}{f_a^2 M_P^2} \left| \left\langle W_{\text{const}} \left[K^i (K^{-1})_i{}^j W_{0j}^* - 3W_0^* \right] \right\rangle \right|, \quad (2.18)$$

where f_a is the decay constant of the R -axion, which is given by

$$f_a = \left[2K_i{}^j r_i r_j \langle \phi_i \rangle \langle \phi_j \rangle \right]^{1/2}, \quad (2.19)$$

with r_i the R charge of the field ϕ_i , and

$$K^i \equiv \frac{\partial K}{\partial \phi_i}, \quad K_i{}^j \equiv \frac{\partial^2 K}{\partial \phi^{*i} \partial \phi_j}, \quad W_{0j}^* \equiv \frac{\partial W_0^*}{\partial \phi^{*j}}. \quad (2.20)$$

⁶We assume that this is the only term that violates the $\text{U}(1)_R$ symmetry; in particular, $\text{U}(1)_R$ -breaking non-renormalizable operators are supposed to be absent, though such operators may be induced if the $\text{U}(1)_R$ symmetry is violated by quantum gravity effect.

Suppose that the VEVs of the $U(1)_R$ charged fields in the hidden sector are much smaller than $\langle S \rangle$. Then, we have

$$f_a \simeq \frac{\sqrt{2}}{9} \langle S \rangle = 5 \times 10^{16} \times \left(\frac{\langle S \rangle}{3 \times 10^{17} \text{ GeV}} \right) \text{ GeV} . \quad (2.21)$$

To estimate the size of the right-hand side in eq. (2.18), we first note that on the assumption that the soft mass scale m_{SUSY} is generated via the gravity mediation effect from the hidden sector, the size of the F -term in the hidden sector, F_Z , is approximately given by $F_Z \sim m_{\text{SUSY}} M_P$, and thus the typical energy scale of the hidden sector physics is $\Lambda_{\text{hid}} \simeq \sqrt{F_Z} \simeq \sqrt{m_{\text{SUSY}} M_P}$. As a result, the hidden-sector contribution to the terms in the square brackets in eq. (2.18) is order

$$\Lambda_{\text{hid}}^3 = (m_{\text{SUSY}} M_P)^{3/2} \simeq 4 \times 10^{33} \times \left(\frac{m_{\text{SUSY}}}{10^4 \text{ GeV}} \right)^{\frac{3}{2}} \text{ GeV}^3 . \quad (2.22)$$

On the other hand, it is found that the contribution from W_{HS} can well be approximated by that of the last term in eq. (2.7) and is given by

$$\left| \left\langle S \frac{\partial W_{HS}^*}{\partial S^*} - 3W_{HS}^* \right\rangle \right| \simeq 5 \times 10^{38} \times \left(\frac{\langle S \rangle}{3 \times 10^{17} \text{ GeV}} \right)^{18} \text{ GeV}^3 , \quad (2.23)$$

for $\Lambda_{HS} = M_P$, which is much larger than that in eq. (2.22).⁷ Finally, W_{const} is estimated as

$$W_{\text{const}} \simeq |\langle W \rangle| = m_{3/2} M_P^2 = 6 \times 10^{40} \times \left(\frac{m_{3/2}}{10^4 \text{ GeV}} \right) \text{ GeV}^3 , \quad (2.24)$$

where $m_{3/2}$ is the gravitino mass. All in all, the R -axion mass m_a is evaluated as

$$m_a \simeq 1 \times 10^5 \times \left(\frac{\langle S \rangle}{3 \times 10^{17} \text{ GeV}} \right)^8 \left(\frac{m_{3/2}}{10^4 \text{ GeV}} \right)^{1/2} \text{ GeV} . \quad (2.25)$$

We see that R -axion is as heavy as other SUSY particles in this model.

As we see above, the bilinear terms $H\bar{H}$ and $h\bar{h}$ are forbidden by the $U(1)_R$ symmetry in this model. The former is effectively induced from W_{HS} after S , H , and $\bar{H}$ develop VEVs, with its coefficient, μ_H , being $\mathcal{O}(m_{\text{SUSY}})$. The generation of the μ -term for $h\bar{h}$ depends on the $U(1)_R$ charges of the hidden sector fields. With the superpotential terms for the hidden sector having $U(1)_R$ charge +2, we expect there are $U(1)_R$ charged fields in the hidden sector. If the SUSY-breaking field Z has $U(1)_R$ charge 19/9, for example, the μ term is generated from a Kähler potential term of the form

$$\frac{1}{\Lambda_Z} \int d^4\theta Z^* h \bar{h} + \text{h.c.} , \quad (2.26)$$

with which we have $\mu = F_Z/\Lambda_Z \simeq \mathcal{O}(m_{\text{SUSY}})$ for $\Lambda_Z \simeq M_P$. In general, a Kähler potential term comprised of Z^* , $S^{(*)}$, $(H\bar{H})^{(*)}$, and $h\bar{h}$ can contribute to the μ -term if the operator

⁷We, however, find that $|F_S|, |F_H|, |F_{\bar{H}}| \ll |F_Z|$ due to the large VEVs of S , H , and $\bar{H}$, and thus the SUSY is broken mainly in the hidden sector.

is allowed by the $U(1)_R$ symmetry. Another contribution can be provided by the coupling to the superconformal compensator field Σ [49–52], which has Weyl weight +1 and $U(1)_R$ charge +2/3. With this field, there is a term allowed by the symmetry,

$$\int d^4\theta \frac{\Sigma^*}{\Sigma^2} S^* h \bar{h} + \text{h.c.}, \quad (2.27)$$

which gives $\mu = \mathcal{O}(m_{3/2}\langle S\rangle/M_P)$. The μ term can also be generated at loop level; for example, one loop diagrams with Higgs bosons and a bino/wino yield a μ term that is suppressed by a loop factor compared with the bino/wino mass [53].

Finally, we discuss the masses of right-handed neutrinos. Since right-handed neutrinos are charged under the $SU(5) \times U(1)$ and $U(1)_R$ symmetries, their masses can be generated only after these symmetries are broken. The dominant contribution to the masses comes from the superpotential term W_{neutrino} in eq. (2.6), which gives rise to a mass matrix for right-handed neutrinos of the form

$$(M_R)_{ij} = c_{ij} \frac{\langle S \rangle |\langle \Phi \rangle|^2}{\Lambda_N^2} \\ \simeq c_{ij} \times 3 \times 10^{13} \times \left(\frac{\langle S \rangle}{3 \times 10^{17} \text{ GeV}} \right) \left(\frac{|\langle \Phi \rangle|}{10^{16} \text{ GeV}} \right)^2 \left(\frac{\Lambda_N}{10^{18} \text{ GeV}} \right)^{-2} \text{ GeV}. \quad (2.28)$$

For the first two generations, this size of the right-handed neutrino masses is large enough to explain the light neutrino masses via the seesaw mechanism [39–43], but may be too small for the third generation, as we see in section 3. We, however, note that a relatively small value of Λ_N may be sufficient to make the right-handed neutrino masses large enough. Moreover, the same operators as in eq. (2.6) with a slightly lower effective cut-off scale can easily be generated by additional fields around the Planck scale; for instance, suppose that there are two singlet chiral superfields, ϕ and $\bar{\phi}$, with $U(1)_R$ charges 19/18 and 17/18, respectively, which have the superpotential interactions

$$W_{\text{ex}} = \lambda_6^i F_i \bar{H} \phi + \mu_\phi \phi \bar{\phi} + \frac{1}{2} \kappa_\phi S \bar{\phi} \bar{\phi}. \quad (2.29)$$

Below the mass scale of μ_ϕ , which is taken to be slightly below the Planck scale, these interactions generate an effective operator

$$W'_{\text{neutrino}} = \frac{\lambda_6^i \lambda_6^j \kappa_\phi}{2\mu_\phi^2} S (F_i^{\alpha\beta} \bar{H}_{\alpha\beta}) (F_j^{\gamma\delta} \bar{H}_{\gamma\delta}), \quad (2.30)$$

which can generate a mass of $\mathcal{O}(10^{14})$ GeV for the third generation right-handed neutrino if $\mu_\phi/M_P = \mathcal{O}(0.1)$.

2.3 Gauge coupling unification

The gauge coupling constants of the $SU(5) \times U(1)$ interactions, g_5 and g_X , are matched onto the SM gauge couplings at the unification scale, $M_{\text{GUT}} \simeq \langle \Phi \rangle$. At tree level, the matching conditions are

$$g_2(M_{\text{GUT}}) = g_3(M_{\text{GUT}}) = g_5(M_{\text{GUT}}), \quad (2.31)$$

and

$$\frac{25}{g_1^2(M_{\text{GUT}})} = \frac{1}{g_5^2(M_{\text{GUT}})} + \frac{24}{g_X^2(M_{\text{GUT}})} . \quad (2.32)$$

We see that g_1 is not necessarily equal to g_2 and g_3 at the unification scale, contrary to the standard SU(5). We, however, note that if we require $g_5(M_{\text{GUT}}) = g_X(M_{\text{GUT}})$, which is expected if the $\text{SU}(5) \times \text{U}(1)$ gauge group is embedded into a simple group such as $\text{SO}(10)$,⁸ then we need $g_1(M_{\text{GUT}}) = g_2(M_{\text{GUT}}) = g_3(M_{\text{GUT}})$ as in the standard SU(5).

Note that in the present scenario M_{GUT} is different from that obtained in the MSSM, $\simeq 2 \times 10^{16}$ GeV, due to the existence of a light color-triplet Higgs multiplet at an intermediate scale. To see its effect on M_{GUT} , let us investigate its contribution to the gauge coupling beta functions. The one-loop beta-function coefficients for the gauge couplings in the MSSM with and without the color-triplet Higgs multiplet, $b_a^{\text{MSSM}+H_C}$ and b_a^{MSSM} ($a = 1, 2, 3$), respectively, are

$$b_a^{\text{MSSM}+H_C} = \begin{pmatrix} 7 \\ 1 \\ -2 \end{pmatrix} \quad \text{and} \quad b_a^{\text{MSSM}} = \begin{pmatrix} 33/5 \\ 1 \\ -3 \end{pmatrix} . \quad (2.33)$$

Since the unification scale is determined by the condition (2.31), the difference in $b_a^{\text{MSSM}+H_C}$ and b_a^{MSSM} for $a = 2, 3$ is relevant. As we see in eq. (2.33), the light color-triplet Higgs multiplet increases the $\text{SU}(3)_C$ beta function above M_{H_C} , while it does not change the $\text{SU}(2)_L$ one. As a result, the unification point becomes higher than that in the MSSM, $\simeq 2 \times 10^{16}$ GeV, which may have some tension with the value of $|\langle \Phi \rangle|$ shown in figure 1, as $M_{\text{GUT}} > |\langle \Phi \rangle|$. The mismatch between these two scales may indicate that there are sizable threshold corrections at the scale M_{GUT} and thus the theory above this scale has a more complicated structure than our model, which is the case if the model is a part of the more fundamental theory based on a larger simple gauge group. A concrete model building for such a theory will be performed on another occasion [46].

2.4 R -parity violation

The $\text{SU}(5) \times \text{U}(1)$ gauge symmetry by itself allows the renormalizable operators

$$W_{R\text{-odd}} = \mu_{FH}^i F_i^{\alpha\beta} \bar{H}_{\alpha\beta} + \kappa_F^i \epsilon_{\alpha\beta\gamma\delta\epsilon} F_i^{\alpha\beta} H^{\gamma\delta} h^\epsilon + \kappa_{\bar{f}}^i \bar{f}_{i\alpha} H^{\alpha\beta} \bar{h}_\beta , \quad (2.34)$$

which yield R -parity violating operators at low energies. All of these operators are eliminated by the $\text{U}(1)_R$ symmetry in our model. At non-renormalizable level, there are operators allowed by the symmetry; for instance, we have

$$\begin{aligned} W_{R\text{-odd}} = & \frac{c_{FH}^i}{\Lambda^6} (H\bar{H}) S^5 F_i^{\alpha\beta} \bar{H}_{\alpha\beta} + \frac{c_F^i}{2\Lambda^4} \epsilon_{\alpha\beta\gamma\delta\epsilon} S^4 F_i^{\alpha\beta} H^{\gamma\delta} h^\epsilon + \frac{\sqrt{2}c_{\bar{f}}^i}{\Lambda^4} S^4 \bar{f}_{i\alpha} H^{\alpha\beta} \bar{h}_\beta \\ & + \frac{\sqrt{2}c_1^{ijk}}{4\Lambda^6} \epsilon_{\alpha\beta\gamma\delta\epsilon} S^5 F_i^{\alpha\beta} F_j^{\gamma\delta} H^{\epsilon\tau} \bar{f}_{k\tau} + \frac{c_2^{ijk}}{\sqrt{2}\Lambda^6} \epsilon_{\alpha\beta\gamma\delta\epsilon} S^5 F_i^{\alpha\beta} H^{\gamma\delta} F_j^{\epsilon\tau} \bar{f}_{k\tau} \\ & + \frac{\sqrt{2}c_3^{ijk}}{\Lambda^6} S^5 \bar{f}_{i\alpha} \bar{f}_{j\beta} \ell_k^c H^{\alpha\beta} + \dots , \end{aligned} \quad (2.35)$$

⁸We have chosen the normalization of the U(1) charge such that the equation $g_5 = g_X$ holds in this case.

where Λ is a cut-off scale and the dots denote operators whose contribution is less significant. At low energies, these operators lead to

$$\begin{aligned}
 W_{R\text{-odd}}^{(\text{eff})} = & \frac{c_{FH}^i \langle \Phi \rangle^3 \langle S \rangle^5}{\Lambda^6} \nu_i^c + \frac{c_F^i \langle \Phi \rangle \langle S \rangle^4}{\Lambda^4} d_i^c D + \frac{c_{\bar{f}}^i \langle \Phi \rangle \langle S \rangle^4}{\Lambda^4} L_i \cdot H_u \\
 & - \frac{(c_1^{ijk} - c_2^{ijk}) \langle \Phi \rangle \langle S \rangle^5}{\Lambda^6} d_i^c Q_j \cdot L_k - \frac{c_2^{ijk} \langle \Phi \rangle \langle S \rangle^5}{\Lambda^6} \epsilon^{abc} d_{ia}^c d_{jb}^c u_{kc}^c \\
 & + \frac{c_3^{ijk} \langle \Phi \rangle \langle S \rangle^5}{\Lambda^6} L_i \cdot L_j e_k^c + \dots,
 \end{aligned} \tag{2.36}$$

where a, b, c denote the color index, ϵ^{abc} is the totally antisymmetric tensor, and $A \cdot B \equiv \epsilon_{mn} A^m B^n$ with m, n the $\text{SU}(2)_L$ index. Although these operators are suppressed by a large power of the cut-off scale Λ , some of them are actually incompatible with the limits on baryon/lepton-number violation [54–56]. Among others, the bound on proton lifetime gives the most stringent constraint; for $\lambda'_{ijk} L_i Q_j d_k^c$ and $(\lambda''_{ijk}/2) u_i^c d_j^c d_k^c$, we have, for instance [56]

$$|\lambda'_{l1k} \lambda''_{11k}| \lesssim 10^{-23} \times \left(\frac{m_{\tilde{d}_k^c}}{10 \text{ TeV}} \right)^2, \tag{2.37}$$

which strongly disfavors the operators in the second line in eq. (2.35). The bound from the washout of baryon asymmetry also restricts the R -parity violating operators [57–61]. For example, we have [60]

$$\frac{c_{\bar{f}}^i \langle \Phi \rangle \langle S \rangle^4}{\Lambda^4} \lesssim \mathcal{O}(10^{-6}) \times |\mu|, \tag{2.38}$$

which excludes the possibility that $c_{\bar{f}}^i = \mathcal{O}(1)$ even for $\Lambda = M_P$.

Considering these severe limits, we introduce R -parity as an extra symmetry, under which the fields in our model transform as⁹

$$\begin{cases} \text{even :} & H, \bar{H}, h, \bar{h}, S \\ \text{odd :} & F_i, \bar{f}_i, \ell_i^c \end{cases}. \tag{2.39}$$

With this assignment, all of the operators in eq. (2.35) are prohibited, while the operators in eqs. (2.4)–(2.7) are still allowed.

3 Flavor structure

In this section, we examine the structure of the Yukawa terms in eq. (2.4) and the neutrino mass matrix, following the discussion in ref. [62]. We define the effective mass matrix for the field F_i^{45} below the symmetry breaking scale as

$$W_R = M_R^{ij} F_i^{45} F_j^{45}. \tag{3.1}$$

As discussed in section 2.2, we assume that it is generated by W_{neutrino} in eq. (2.6) or some additional effect around the Planck scale, and that the values of M_R^{ij} can be as large as $\mathcal{O}(10^{14}) \text{ GeV}$.

⁹When we consider the singlet fields in eq. (2.29), we assign R -parity odd to ϕ and $\bar{\phi}$.

We adopt the basis where λ_2^{ij} is real and diagonal, without loss of generality: $\lambda_2^{ij} = \lambda_{2,D}^i \delta^{ij}$. In this basis, λ_1^{ij} is a generic complex symmetric matrix and can be diagonalized with a unitary matrix, which we denote by U_1 :

$$\lambda_1 = U_1^T \lambda_{1,D} U_1, \quad (3.2)$$

where $\lambda_{1,D}$ is a diagonal matrix whose components are real and positive. Now it is convenient to separate out the phase factors in the matrix U_1 as many as possible, as done in ref. [63]: $(U_1)_{ij} = e^{i\phi_{ij}} v_{ij}$, with ϕ_{ij} and v_{ij} both real (summation over the indices i, j is not taken). We then define the following diagonal phase matrices:

$$P_1 = \text{diag} \left(e^{i\phi_{11}}, e^{i\phi_{21}}, e^{i\phi_{31}} \right), \quad (3.3)$$

$$P_2 = \text{diag} \left(1, e^{i(\phi_{12}-\phi_{11})}, e^{i(\phi_{13}-\phi_{11})} \right). \quad (3.4)$$

By means of these matrices, we decompose the matrix U_1 as

$$U_1 = P_1 V^\dagger P_2, \quad (3.5)$$

to eliminate the phases in the first row and column in U_1 . We further define $P \equiv P_1^2 \exp(-i2 \sum_i \phi_{i1}/3)$, for which $\det P = 1$. The phase factors P_2 and $\exp(-i \sum_i \phi_{i1}/3)$ can be absorbed by the field redefinition without changing the λ_2 couplings:

$$F \rightarrow \exp \left(-i \sum_i \phi_{i1}/3 \right) P_2^* F, \quad \bar{f} \rightarrow \exp \left(i \sum_i \phi_{i1}/3 \right) P_2 \bar{f}. \quad (3.6)$$

Since U_1 is a 3×3 unitary matrix, it has nine free parameters. P_1 and P_2 have five phases in total, and thus V has four free parameters. With the condition $\det P = 1$, the phase matrix P has two degrees of freedom.

In the above basis, the λ_1 matrix is written as

$$\lambda_1 = V^* P \lambda_{1,D} V^\dagger. \quad (3.7)$$

After the field redefinition, the coupling matrix λ_3 has been transformed as

$$\lambda_3 \rightarrow \exp \left(i \sum_i \phi_{i1}/3 \right) P_2 \lambda_3. \quad (3.8)$$

We absorb the factor $\exp(i \sum_i \phi_{i1}/3)$ with the field redefinition

$$\ell^c \rightarrow \exp \left(-i \sum_i \phi_{i1}/3 \right) \ell^c. \quad (3.9)$$

The remaining part, $P_2 \lambda_3$, is diagonalized with unitary matrices U_ℓ and U_{ℓ^c} as

$$\lambda_{3,D} = U_\ell^\dagger P_2 \lambda_3 U_{\ell^c}. \quad (3.10)$$

Finally, the matrix M_R^{ij} in eq. (3.1) has been transformed as

$$M_R \rightarrow \exp \left(-2i \sum_i \phi_{i1}/3 \right) P_2^* M_R P_2^* \equiv M'_R, \quad (3.11)$$

which is diagonalized as

$$M_{R,D} = U_{\nu^c}^T M'_R U_{\nu^c}, \quad (3.12)$$

where $M_{R,D} \equiv \text{diag}(M_{R_1}, M_{R_2}, M_{R_3})$ with M_{R_i} real and positive.

After these field redefinitions, the terms in W_{Yukawa} and W_R are written as

$$\begin{aligned} W_{\text{eff}} = & \frac{1}{4} (V^* P \lambda_{1,D} V^\dagger)^{ij} \epsilon_{\alpha\beta\gamma\delta\epsilon} F_i^{\alpha\beta} F_j^{\gamma\delta} h^\epsilon + \sqrt{2} \lambda_{2,D}^i F_i^{\alpha\beta} \bar{f}_{i\alpha} \bar{h}_\beta \\ & + (U_\ell \lambda_{3,D} U_{\ell^c}^\dagger)^{ij} \bar{f}_{i\alpha} \ell_j^c h^\alpha + (U_{\nu^c}^* M_{R,D} U_{\nu^c}^\dagger)^{ij} F_i^{45} F_j^{45}. \end{aligned} \quad (3.13)$$

Now we express the fields in the above equation in terms of the component fields defined as follows:

$$\begin{aligned} F_i^{ab} &= \frac{1}{\sqrt{2}} \epsilon^{abc} V_{ij} P_{jj}^* d_{jc}^c, & F_i^{an} &= \frac{1}{\sqrt{2}} Q_i^{an}, & F_i^{45} &= \frac{1}{\sqrt{2}} U_{\nu^c}^{ij} \nu_j^c, \\ \bar{f}_{ia} &= u_{ia}^c, & \bar{f}_{in} &= \epsilon_{nm} U_\ell^{*ij} L_{jm}, & \ell_i^c &= U_{\ell^c}^{ij} e_j^c \end{aligned} \quad (3.14)$$

where $m, n = 4, 5$ denote the $\text{SU}(2)_L$ index, and

$$h = \begin{pmatrix} D^1 \\ D^2 \\ D^3 \\ H_d^0 \\ H_d^- \end{pmatrix}, \quad \bar{h} = \begin{pmatrix} \bar{D}_1 \\ \bar{D}_2 \\ \bar{D}_3 \\ H_u^0 \\ -H_u^+ \end{pmatrix}. \quad (3.15)$$

We then have

$$\begin{aligned} W_{\text{eff}} = & (\lambda_{1,D} V^\dagger)^{ij} d_i^c (Q_j \cdot H_d) + \lambda_{2,D}^i u_i^c (Q_i \cdot H_u) + \lambda_{3,D}^i (H_d \cdot L_i) e_i^c \\ & + (U_{\nu^c}^T \lambda_{2,D} U_\ell^*)^{ij} \nu_i^c (L_j \cdot H_u) + \frac{1}{2} M_{R_i} \nu_i^c \nu_i^c \\ & - \frac{1}{2} (V^* P \lambda_{1,D} V^\dagger)^{ij} \epsilon_{abc} (Q_i^a \cdot Q_j^b) D^c + (\lambda_{1,D} V^\dagger U_{\nu^c})^{ij} d_i^c \nu_j^c D \\ & + (P^* V^T \lambda_{2,D})^{ij} \epsilon^{abc} d_{ia}^c u_{jb}^c \bar{D}_c + (U_\ell^\dagger \lambda_{2,D})^{ij} (L_i \cdot Q_j) \bar{D} + (U_\ell \lambda_{3,D})^{ij} u_i^c e_j^c D. \end{aligned} \quad (3.16)$$

From the first two terms, we see that the up- and down-type quark masses are diagonalized if we define the components of the left-handed quark fields as

$$Q_i = \begin{pmatrix} u_i \\ V_{ij} d_j \end{pmatrix}. \quad (3.17)$$

Therefore, the matrix V is equal to the Cabibbo-Kobayashi-Maskawa (CKM) matrix. The eigenvalues of $\lambda_{1,D}$, $\lambda_{2,D}$, and $\lambda_{3,D}$ correspond to the down-type quark, up-type quark, and charged lepton Yukawa couplings, respectively.

We here note in passing that the down-type quark and charged-lepton Yukawa couplings do not necessarily unify in flipped $\text{SU}(5)$ GUT models, contrary to the standard $\text{SU}(5)$. This feature could be advantageous given that the unification of Yukawa couplings is generically imperfect in the MSSM, especially for the first two generations.

At the mass scale of right-handed neutrinos, we integrate out these fields to obtain the dimension-five operator for the light-neutrino mass matrix [39–43]:

$$W_\nu = -\frac{1}{2} \left(U_\ell^\dagger \lambda_{2,D} U_{\nu^c} M_{R,D}^{-1} U_{\nu^c}^T \lambda_{2,D} U_\ell^* \right)^{ij} (L_i \cdot H_u)(L_j \cdot H_u) . \quad (3.18)$$

For convenience, we define the neutrino mass matrix as

$$m_\nu^{ij} \equiv - \left(\lambda_{2,D} U_{\nu^c} M_{R,D}^{-1} U_{\nu^c}^T \lambda_{2,D} \right)^{ij} \langle H_u^0 \rangle^2 . \quad (3.19)$$

This mass matrix can be diagonalized with a unitary matrix:

$$m_\nu^D = U_\nu^T m_\nu U_\nu . \quad (3.20)$$

We express the left-handed lepton fields in terms of the mass eigenstates as

$$L_i = \begin{pmatrix} U_{ij} \nu_j \\ e_i \end{pmatrix} , \quad (3.21)$$

where U is the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix. We then find that it is related to the unitary matrices U_ℓ and U_ν as

$$U = U_\ell^T U_\nu . \quad (3.22)$$

Now recall that the couplings $\lambda_{2,D}^i$ correspond to the up-type quark Yukawa couplings. The eigenvalues of the mass matrix m_ν are, therefore, approximately given by

$$m_{\nu_i} \simeq \frac{m_{q_i}^2}{M_N} , \quad (3.23)$$

where m_{q_i} are quark masses and M_N represents the scale of right-handed neutrino masses. By requiring $m_{\nu_3} \lesssim 0.1 \text{ eV}$, we obtain a rough lower limit on the right-handed neutrino mass scale: $M_N \gtrsim 3 \times 10^{14} \text{ GeV}$. The relation (3.19) also suggests that the components in m_ν are hierarchical unless $M_{R,D}$ is hierarchical, and thus the matrix U_ν is close to the unity. In this case, the matrix U_ℓ can well be approximated by the PMNS matrix as

$$U_\ell \simeq U^T . \quad (3.24)$$

We focus on this case in the following discussion.

4 Proton decay

In our model, proton decay is caused by the exchange of the color-triplet Higgs multiplets or the SU(5) gauge bosons; the former induces both the dimension-five and dimension-six operators, while the latter generates only the dimension-six operators.

4.1 Dimension-five proton decay

Let us first discuss the dimension-five proton decay operators induced by the color-triplet Higgs exchange. In many SUSY GUT models, such as the minimal SUSY SU(5) [13, 14], these operators give the dominant contribution [64, 65], with the main decay mode being $p \rightarrow K^+ \bar{\nu}$ [66]. In the standard flipped SU(5), however, this contribution is extremely small since the μ -term for $H\bar{H}$, μ_H , is $\mathcal{O}(m_{\text{SUSY}})$ and thus the coefficients of the dimension-five operators are suppressed by a factor of $\mathcal{O}(m_{\text{SUSY}}/M_{\text{GUT}}^2)$ due to the chirality flip. In our model, the suppression of the dimension-five proton decay by a small μ_H also exists, whereas one of the color-triplet Higgs multiplets has a mass much smaller than M_{GUT} . It is, therefore, worth checking if the dimension-five proton decay rate is still small enough to evade the current experimental limits.

The exchange of the color-triplet Higgs multiplets is induced by the interactions in eq. (3.16). By integrating out the fields D , $\bar{D}$, d_H^c , and $d_{\bar{H}}^c$, we obtain the following dimension-five superpotential-type operators that violate baryon and lepton numbers:

$$W_5^{\text{eff}} = \frac{\mu_H}{M_{H_C} M_{\bar{H}_C}} \left[\frac{1}{2} (V^* P \lambda_{1,D} V^\dagger)^{ij} (\lambda_{2,D} U_\ell^*)^{kl} \epsilon_{abc} (Q_i^a \cdot Q_j^b) (Q_k^c \cdot L_l) \right. \\ \left. - (\lambda_{2,D} V P^*)^{ij} (U_\ell \lambda_{3,D})^{kl} \epsilon^{abc} u_{ia}^c d_{jb}^c u_{kc}^c e_l^c \right]. \quad (4.1)$$

As we see, these operators receive a suppression factor μ_H/M_{H_C} in addition to the ordinary factor $M_{\bar{H}_C}^{-1} \simeq M_{\text{GUT}}^{-1}$ for the dimension-five operators induced at the GUT scale. To see if this suppression factor is sufficient to evade the current experimental bound, we note that in the minimal SUSY SU(5), the lifetime of $p \rightarrow K^+ \bar{\nu}$ is predicted to be $\tau(p \rightarrow K^+ \bar{\nu}) = \mathcal{O}(10^{30-31})$ years for the TeV-scale SUSY particles, the color-triplet Higgs mass of $\simeq 10^{16}$ GeV, and $\tan \beta = \mathcal{O}(1)$ [67, 68]. Given that for $\mu_H \simeq 1$ TeV and $M_{H_C} \simeq 10^{11}$ GeV the suppression factor is $\mu_H/M_{H_C} = 10^{-8}$, we find that the proton lifetime in our model is much longer than the current experimental bound, $\tau(p \rightarrow K^+ \bar{\nu}) > 6.6 \times 10^{33}$ years [69, 70], for the same parameter choice.¹⁰ If m_{SUSY} is much larger than the TeV scale, the suppression factor μ_H/M_{H_C} can be moderate. In this case, however, the dimension-five proton decay evades the experimental bound even in the minimal SUSY SU(5) [71–82], and thus there is no constraint on the operators in eq. (4.1).

A potentially dangerous contribution to the dimension-five proton decay in our model is provided by the cut-off suppressed operators. We find that the following operators are allowed by the symmetries of the theory:

$$W'_5 = \frac{c_1^{ijkl}}{\Lambda^2} \epsilon_{\alpha\beta\gamma\delta\epsilon} S F_i^{\alpha\beta} F_j^{\gamma\delta} F_k^{\epsilon\tau} \bar{f}_{l\tau} + \frac{c_2^{ijkl}}{\Lambda^2} S F_i^{\alpha\beta} \bar{f}_{j\alpha} \bar{f}_{k\beta} \ell_l^c. \quad (4.2)$$

It turns out that if c_1^{ijkl} or c_2^{ijkl} is $\mathcal{O}(1)$, then the rate of the dimension-five proton decay is too large for $m_{\text{SUSY}} \lesssim \mathcal{O}(10^8)$ GeV, even for $\Lambda = M_P$ [83]; for a smaller m_{SUSY} , the

¹⁰Strictly speaking, this estimate is not accurate since the coefficients of the dimension-five operators in the minimal SUSY SU(5) have different flavor structure compared with those in eq. (4.1). We, however, find that the coefficients are not so much different in these two cases, and thus our conclusion is unchanged.

limits on these coefficients from the proton decay bound get more severe, decreasing as $\propto m_{\text{SUSY}}$. We, however, note that it is quite likely that the coefficients c_1^{ijkl} and c_2^{ijkl} are suppressed by the same mechanism as that explains the hierarchical structure of the SM quark and lepton Yukawa couplings. Indeed, if we consider a flavor-dependent $U(1)_R$ charge assignment to account for the Yukawa structure, then the coefficients c_1^{ijkl} and c_2^{ijkl} can be highly suppressed for $i, j, k, l = 1, 2$, with which we can evade the proton decay bound. In this paper, we just assume that a certain mechanism makes these coefficients small enough to evade the experimental limit, as often done in the literature.

4.2 Dimension-six proton decay

The suppression of the dimension-five proton decay operators (4.1) by a small μ_H is due to the chirality flip that is required to generate the operators. The exchange of the light color-triplet Higgs multiplet induces dimension-six operators as well via chirality-conserving diagrams. Since these dimension-six operators do not suffer from the suppression by μ_H , they may dominate the contribution from W_5^{eff} even if they have larger mass dimensions than those in eq. (4.1). In addition, the exchange of the $SU(5)$ gauge bosons also induces the dimension-six proton decay operators, which gives the dominant contribution to proton decay in the standard flipped $SU(5)$ model [84–88]. We now study the dimension-six proton decay in our model, taking account of both the light color-triplet Higgs multiplet and $SU(5)$ gauge boson exchange processes.

The dimension-six proton decay operators are expressed as

$$\mathcal{L}_6^{\text{eff}} = C_{6(1)}^{ijkl} \mathcal{O}_{ijkl}^{6(1)} + C_{6(2)}^{ijkl} \mathcal{O}_{ijkl}^{6(2)}, \quad (4.3)$$

where

$$\mathcal{O}_{ijkl}^{6(1)} = \int d^2\theta d^2\bar{\theta} \epsilon_{abc} \epsilon_{mn} (u_i^{c\dagger})^a (d_j^{c\dagger})^b e^{-\frac{2}{3}g'B} (e^{2g_3G} Q_k^m)^c L_l^n, \quad (4.4)$$

$$\mathcal{O}_{ijkl}^{6(2)} = \int d^2\theta d^2\bar{\theta} \epsilon_{abc} \epsilon_{mn} Q_i^{am} Q_j^{bn} e^{\frac{2}{3}g'B} (e^{-2g_3G} u_k^{c\dagger})^c e_l^{c\dagger}, \quad (4.5)$$

with G and B the $SU(3)_C$ and $U(1)_Y$ gauge vector superfields, respectively.

First, at the scale M_{V_5} , we integrate out the $SU(5)$ gauge bosons to obtain these proton decay operators. The relevant gauge interaction terms are

$$K_{\text{gauge}} = \sqrt{2}g_5 (-\epsilon_{mn} (u_a^c)^\dagger X_a^m U_\ell^* L^n + \epsilon^{abc} (Q^{am})^\dagger X_b^m V P^* d_c^c + \epsilon_{mn} (\nu^c)^\dagger U_{\nu^c}^\dagger X_a^m Q^{an} + \text{h.c.}), \quad (4.6)$$

where X_a^m denote the $SU(5)$ gauge vector superfields. The tree-level exchange of the $SU(5)$ gauge bosons then yields

$$C_{6(1)}^{ijkl}(M_{V_5}) = \frac{g_5^2}{M_{V_5}^2} (U_\ell^*)_{il} V_{kj}^* e^{i\varphi_j} = \frac{1}{|\langle\Phi\rangle|^2} (U_\ell^*)_{il} V_{kj}^* e^{i\varphi_j},$$

$$C_{6(2)}^{ijkl}(M_{V_5}) = 0, \quad (4.7)$$

where φ_i are the phases in the phase matrix P , defined by $P_{ij} \equiv e^{i\varphi_i} \delta_{ij}$. The contribution of the exchange of the heavier color-triplet Higgs multiplet is much smaller than the one

in eq. (4.7) due to the suppression by small Yukawa couplings. We therefore neglect this contribution in the following analysis.

These Wilson coefficients are evolved down to the mass scale of M_{H_C} using the renormalization group equations (RGEs). The renormalization factor can be computed using the results in ref. [89] at one-loop level¹¹ as

$$C_{6(1)}^{ijkl}(M_{H_C}) = A_1^H C_{6(1)}^{ijkl}(M_{V_5}), \quad (4.8)$$

with

$$A_1^H = \left[\frac{\alpha_3(M_{H_C})}{\alpha_3(M_{V_5})} \right]^{\frac{2}{3}} \left[\frac{\alpha_2(M_{H_C})}{\alpha_2(M_{V_5})} \right]^{-\frac{3}{2}} \left[\frac{\alpha_1(M_{H_C})}{\alpha_1(M_{V_5})} \right]^{-\frac{11}{210}}. \quad (4.9)$$

At this scale, we integrate out the light color-triplet Higgs at tree level, which contributes only to the coefficient $C_{6(2)}^{ijkl}$:

$$C_{6(2)}^{ijkl}(M_{H_C}) = -\frac{1}{2M_{H_C}^2} (V^* P \lambda_{1,D} V^\dagger)^{ij} (U_\ell^* \lambda_{3,D})^{kl}. \quad (4.10)$$

We then run both of these coefficients down to the hadronic scale. The resultant renormalization factors are given as

$$\begin{aligned} C_{6(1)}^{ijkl}(\mu_{\text{had}}) &= A_1^R C_{6(1)}^{ijkl}(M_{H_C}), \\ C_{6(2)}^{ijkl}(\mu_{\text{had}}) &= A_2^R C_{6(2)}^{ijkl}(M_{H_C}), \end{aligned} \quad (4.11)$$

with

$$\begin{aligned} A_1^R &= \left[\frac{\alpha_3(m_{\text{SUSY}})}{\alpha_3(M_{H_C})} \right]^{\frac{4}{9}} \left[\frac{\alpha_2(m_{\text{SUSY}})}{\alpha_2(M_{H_C})} \right]^{-\frac{3}{2}} \left[\frac{\alpha_1(m_{\text{SUSY}})}{\alpha_1(M_{H_C})} \right]^{-\frac{1}{18}} \\ &\quad \times \left[\frac{\alpha_3(m_Z)}{\alpha_3(m_{\text{SUSY}})} \right]^{\frac{2}{7}} \left[\frac{\alpha_2(m_Z)}{\alpha_2(m_{\text{SUSY}})} \right]^{\frac{27}{38}} \left[\frac{\alpha_1(m_Z)}{\alpha_1(m_{\text{SUSY}})} \right]^{-\frac{11}{82}} A_L, \\ A_2^R &= \left[\frac{\alpha_3(m_{\text{SUSY}})}{\alpha_3(M_{H_C})} \right]^{\frac{4}{9}} \left[\frac{\alpha_2(m_{\text{SUSY}})}{\alpha_2(M_{H_C})} \right]^{-\frac{3}{2}} \left[\frac{\alpha_1(m_{\text{SUSY}})}{\alpha_1(M_{H_C})} \right]^{-\frac{23}{198}} \\ &\quad \times \left[\frac{\alpha_3(m_Z)}{\alpha_3(m_{\text{SUSY}})} \right]^{\frac{2}{7}} \left[\frac{\alpha_2(m_Z)}{\alpha_2(m_{\text{SUSY}})} \right]^{\frac{27}{38}} \left[\frac{\alpha_1(m_Z)}{\alpha_1(m_{\text{SUSY}})} \right]^{-\frac{23}{82}} A_L, \end{aligned} \quad (4.12)$$

where μ_{had} is the hadronic scale, for which we use $\mu_{\text{had}} = 2 \text{ GeV}$. The renormalization factors below the SUSY-breaking scale are computed from the results in ref. [91]. A_L is the long-distance QCD renormalization factor, which is computed in ref. [92] at the two-loop level: $A_L = 1.24$.

In the standard flipped SU(5), $p \rightarrow \pi^0 e^+ / \mu^+$ are the main decay channels. The relevant effective operators for these modes are

$$\mathcal{L}(p \rightarrow \pi^0 \ell_i^+) = C_{RL}(ud\ell_i)[\epsilon_{abc}(u_R^a d_R^b)(u_L^c \ell_{Li})] + C_{LR}(ud\ell_i)[\epsilon_{abc}(u_L^a d_L^b)(u_R^c \ell_{Ri})], \quad (4.13)$$

¹¹The two-loop RGEs for these coefficients are also available [90].

where

$$\begin{aligned} C_{RL}(ud\ell_i) &= C_{6(1)}^{11i}, \\ C_{LR}(ud\ell_i) &= V_{j1}[C_{6(2)}^{1j1i} + C_{6(2)}^{j11i}]. \end{aligned} \quad (4.14)$$

The partial decay width can then be computed as

$$\Gamma(p \rightarrow \pi^0 \ell_i^+) = \frac{m_p}{32\pi} \left(1 - \frac{m_\pi^2}{m_p^2}\right)^2 [|\mathcal{A}_L(p \rightarrow \pi^0 \ell_i^+)|^2 + |\mathcal{A}_R(p \rightarrow \pi^0 \ell_i^+)|^2], \quad (4.15)$$

where we have neglected the lepton mass for simplicity, m_p and m_π denote the masses of the proton and pion, respectively, and

$$\begin{aligned} \mathcal{A}_L(p \rightarrow \pi^0 \ell_i^+) &= C_{RL}(ud\ell_i) \langle \pi^0 | (ud)_{RuL} | p \rangle_{\ell_i}, \\ \mathcal{A}_R(p \rightarrow \pi^0 \ell_i^+) &= C_{LR}(ud\ell_i) \langle \pi^0 | (ud)_{RuL} | p \rangle_{\ell_i}. \end{aligned} \quad (4.16)$$

For the matrix elements $\langle \pi^0 | (ud)_{RuL} | p \rangle_{\ell_i}$, we use the values obtained with QCD lattice simulations in ref. [93]:

$$\langle \pi^0 | (ud)_{RuL} | p \rangle_e = -0.131(4)(13) \text{ GeV}^2, \quad (4.17)$$

$$\langle \pi^0 | (ud)_{RuL} | p \rangle_\mu = -0.118(3)(12) \text{ GeV}^2, \quad (4.18)$$

where the first and second parentheses show the statistical and systematic uncertainties, respectively, and the subscript $\ell_i = e, \mu$ shows that the matrix element is computed for the final state with a charged antilepton ℓ_i^+ .

By using the above equations, we can express the partial decay widths of the $p \rightarrow \pi^0 e^+/\mu^+$ decay modes as

$$\begin{aligned} \Gamma(p \rightarrow \pi^0 e^+) &= \frac{m_p}{32\pi} \left(1 - \frac{m_\pi^2}{m_p^2}\right)^2 |V_{ud}|^2 |(U_\ell)_{11}|^2 \left(\langle \pi^0 | (ud)_{RuL} | p \rangle_e \right)^2 \\ &\quad \times \left[\frac{(A_1^H)^2 (A_1^R)^2}{|\langle \Phi \rangle|^4} + \frac{(A_2^R)^2}{M_{H_C}^4} \{y_d(M_{H_C}) y_e(M_{H_C})\}^2 \right], \end{aligned} \quad (4.19)$$

$$\begin{aligned} \Gamma(p \rightarrow \pi^0 \mu^+) &= \frac{m_p}{32\pi} \left(1 - \frac{m_\pi^2}{m_p^2}\right)^2 |V_{ud}|^2 |(U_\ell)_{12}|^2 \left(\langle \pi^0 | (ud)_{RuL} | p \rangle_\mu \right)^2 \\ &\quad \times \left[\frac{(A_1^H)^2 (A_1^R)^2}{|\langle \Phi \rangle|^4} + \frac{(A_2^R)^2}{M_{H_C}^4} \{y_d(M_{H_C}) y_\mu(M_{H_C})\}^2 \right], \end{aligned} \quad (4.20)$$

where y_d , y_e , and y_μ are the Yukawa couplings of down quark, electron, and muon, respectively. As we see, these decay widths do not depend on the phases φ_i in P . In addition, the SU(5) gauge boson and color-triplet Higgs multiplet exchange processes do not interfere with each other since the chirality of the charged lepton in the final state is different and thus these two processes are in principle distinguishable.¹²

¹²Strictly speaking, for $p \rightarrow \pi^0 \mu^+$, there is a correction which depends on the interference at $\mathcal{O}(m_\mu/m_p)$. However, its effect is negligible except for a narrow parameter region around $|\mathcal{A}_L| \simeq |\mathcal{A}_R|$, and is at most $2m_\mu/m_p \simeq 0.2$ even for $|\mathcal{A}_L| = |\mathcal{A}_R|$.

Since the color-triplet Higgs exchange process is induced by the Yukawa interactions, its contribution tends to be more significant for the decay modes that contain the second generations. In addition, this process contributes to only the decay channels that contain a charged lepton. Motivated by this observation, we also consider the $p \rightarrow K^0 e^+/\mu^+$ decay modes. The effective interactions in this case are given by

$$\mathcal{L}(p \rightarrow K^0 \ell_i^+) = C_{RL}(usu\ell_i)[\epsilon_{abc}(u_R^a s_R^b)(u_L^c \ell_{Li})] + C_{LR}(usu\ell_i)[\epsilon_{abc}(u_L^a s_L^b)(u_R^c \ell_{Ri})], \quad (4.21)$$

with

$$\begin{aligned} C_{RL}(usu\ell_i) &= C_{6(1)}^{121i}, \\ C_{LR}(usu\ell_i) &= V_{j2}[C_{6(2)}^{1j1i} + C_{6(2)}^{j11i}]. \end{aligned} \quad (4.22)$$

We then obtain the partial decay width

$$\Gamma(p \rightarrow K^0 \ell_i^+) = \frac{m_p}{32\pi} \left(1 - \frac{m_K^2}{m_p^2}\right)^2 [|\mathcal{A}_L(p \rightarrow K^0 \ell_i^+)|^2 + |\mathcal{A}_R(p \rightarrow K^0 \ell_i^+)|^2], \quad (4.23)$$

where m_K is the kaon mass and

$$\begin{aligned} \mathcal{A}_L(p \rightarrow K^0 \ell_i^+) &= C_{RL}(usu\ell_i) \langle K^0 | (us)_{RuL} | p \rangle_{\ell_i}, \\ \mathcal{A}_R(p \rightarrow K^0 \ell_i^+) &= C_{LR}(usu\ell_i) \langle K^0 | (us)_{RuL} | p \rangle_{\ell_i}. \end{aligned} \quad (4.24)$$

The matrix elements are again computed in ref. [93]:

$$\langle K^0 | (us)_{RuL} | p \rangle_e = 0.103(3)(11) \text{ GeV}^2, \quad (4.25)$$

$$\langle K^0 | (us)_{RuL} | p \rangle_\mu = 0.099(2)(10) \text{ GeV}^2. \quad (4.26)$$

By using these results, we obtain

$$\begin{aligned} \Gamma(p \rightarrow K^0 e^+) &= \frac{m_p}{32\pi} \left(1 - \frac{m_K^2}{m_p^2}\right)^2 |V_{us}|^2 |(U_\ell)_{11}|^2 \left(\langle K^0 | (us)_{RuL} | p \rangle_e \right)^2 \\ &\quad \times \left[\frac{(A_1^H)^2 (A_1^R)^2}{|\langle \Phi \rangle|^4} + \frac{(A_2^R)^2}{M_{H_C}^4} \{y_s(M_{H_C}) y_e(M_{H_C})\}^2 \right], \end{aligned} \quad (4.27)$$

$$\begin{aligned} \Gamma(p \rightarrow K^0 \mu^+) &= \frac{m_p}{32\pi} \left(1 - \frac{m_K^2}{m_p^2}\right)^2 |V_{us}|^2 |(U_\ell)_{12}|^2 \left(\langle K^0 | (us)_{RuL} | p \rangle_\mu \right)^2 \\ &\quad \times \left[\frac{(A_1^H)^2 (A_1^R)^2}{|\langle \Phi \rangle|^4} + \frac{(A_2^R)^2}{M_{H_C}^4} \{y_s(M_{H_C}) y_\mu(M_{H_C})\}^2 \right], \end{aligned} \quad (4.28)$$

where y_s is the strange quark Yukawa coupling. Again, these decay widths are independent of the phases φ_i , and there is no interference between the SU(5) gauge boson and color-triplet Higgs contributions (see, however, footnote 12 for an $\mathcal{O}(m_\mu/m_p)$ correction).

As we see in eqs. (4.19), (4.20), (4.27), and (4.28), the proton decay rates depend on the unitary matrix U_ℓ . Motivated by the discussion in section 3, we set $U_\ell = U^T$

Figure 2. The lifetime of each decay mode as a function of the color-triplet Higgs mass M_{H_C} for $|\langle\Phi\rangle| = 10^{16} \text{ GeV}$ and $m_{\text{SUSY}} = 10^4 \text{ GeV}$. The solid blue and dashed green lines correspond to $\tan\beta \equiv \langle H_u^0 \rangle / \langle H_d^0 \rangle = 3$ and 10, respectively. The gray shades represent the current experimental limits from Super-Kamiokande [97–99], and the horizontal black dotted lines show the sensitivities of Hyper-Kamiokande [100].

with U the PMNS matrix; i.e., $(U_\ell)_{11} = \cos\theta_{12}\cos\theta_{13}$ and $(U_\ell)_{12} = -\sin\theta_{12}\cos\theta_{23} - \cos\theta_{12}\sin\theta_{23}\sin\theta_{13}e^{i\delta}$, for the normal ordering, where $\theta_{12}, \theta_{23}, \theta_{13}$ are the neutrino mixing angles and δ is the neutrino Dirac CP phase. We use the values obtained by the NuFIT 5.0 global analysis of neutrino oscillation measurements for these parameters [94, 95], assuming the normal ordering. For other input parameters, we use the values recommended by the Particle Data Group [96].

In figure 2, we show the lifetime of each decay mode as a function of the color-triplet Higgs mass M_{H_C} for $|\langle\Phi\rangle| = 10^{16} \text{ GeV}$ and $m_{\text{SUSY}} = 10^4 \text{ GeV}$. The solid blue

and dashed green lines correspond to $\tan\beta \equiv \langle H_u^0 \rangle / \langle H_d^0 \rangle = 3$ and 10, respectively. The current experimental limits from Super-Kamiokande, $\tau(p \rightarrow \pi^0 e^+) > 1.6 \times 10^{34}$ years [97], $\tau(p \rightarrow \pi^0 \mu^+) > 7.7 \times 10^{33}$ years [97], $\tau(p \rightarrow K^0 e^+) > 1.0 \times 10^{33}$ years [98], and $\tau(p \rightarrow K^0 \mu^+) > 1.6 \times 10^{33}$ years [99], are shown in the gray shaded area, and the expected 90% CL limits from the 10-year run of Hyper-Kamiokande [100] are shown in the horizontal black dotted lines (if available). As we see, $p \rightarrow \pi^0 \mu^+$ and $p \rightarrow K^0 \mu^+$ are most sensitive to the proton decay induced by the light color-triplet Higgs exchange, setting a lower limit on its mass, $M_{H_C} \gtrsim 5 \times 10^{11}$ GeV for $\tan\beta \gtrsim 3$. The lifetime is shorter for the decay mode that contains a second-generation quark or lepton in the final state due to their larger Yukawa couplings. It is also found that the decay rate increases for a larger $\tan\beta$, since both the down-type quark and charged lepton Yukawa couplings are inversely proportional to $\cos\beta$. For a sufficiently large M_{H_C} , the lifetimes approach asymptotic values, which are determined by the SU(5) gauge boson exchange process.

We note that we can in principle distinguish the proton decay induced by the color-triplet Higgs exchange from that by the SU(5) gauge boson exchange in our model by measuring the chirality of the final-state lepton. As we see in eqs. (4.7) and (4.10), the SU(5) gauge boson generates only the operator that contains a left-handed lepton,¹³ while the color-triplet Higgs exchange induces that with a right-handed lepton. We can thus discriminate these two contributions by investigating the polarization of the antilepton in the decay products, offering an interesting way to test this model.

Our model predicts that if the mass of the lighter color-triplet Higgs field is $\lesssim 10^{12}$ GeV, then proton decay can be discovered at Hyper-Kamiokande in the $p \rightarrow \pi^0 \mu^+$ channel, and presumably in the $p \rightarrow K^0 \mu^+$ channel as well. Notice that without the color-triplet Higgs exchange contribution, i.e., in the case of the standard flipped SU(5), $\tau(p \rightarrow \pi^0 e^+) < \tau(p \rightarrow \pi^0 \mu^+) < \tau(p \rightarrow K^0 \mu^+)$ as can be seen from figure 2. In the minimal SUSY SU(5), on the other hand, the most promising decay channel is either $p \rightarrow K^+ \bar{\nu}$ or $p \rightarrow \pi^0 e^+$, depending on the sfermion mass scale [74, 75, 82]. It is worth noting that $p \rightarrow K^+ \bar{\nu}$ does not occur in the standard flipped SU(5) [62, 88], as well as in our model since the lighter color-triplet Higgs multiplet does not couple to the left-handed neutrinos. In any cases, the $p \rightarrow \pi^0 \mu^+$ and $p \rightarrow K^0 \mu^+$ modes are usually subdominant, and therefore offer a promising way to distinguish our model from others, which highly motivates to search for the $p \rightarrow \pi^0 \mu^+$ and $p \rightarrow K^0 \mu^+$ modes in future proton decay experiments, in addition to the standard channels $p \rightarrow \pi^0 e^+$ and $p \rightarrow K^+ \bar{\nu}$.

5 Conclusion and discussion

We have constructed a flipped SU(5) SUSY GUT model which is invariant under a global $U(1)_R$ symmetry — except the constant term in the superpotential. This $U(1)_R$ symmetry eliminates Higgs potential terms up to sufficiently higher mass dimensions such that the SU(5)-breaking Higgs fields acquire a GUT-scale VEV in the flat direction after SUSY is spontaneously broken. In addition, the μ -terms of the Higgs fields are also forbidden by

¹³This is a specific feature of the flipped SU(5); in the case of the standard SU(5), for instance, the SU(5) gauge bosons contribute to the both types of the operators.

the $U(1)_R$ symmetry and, in particular, the MSSM Higgs superfields naturally have a mass much smaller than the GUT scale through the missing partner mechanism. The masses of right-handed neutrinos in our model are generated by non-renormalizable operators suppressed by a cut-off scale, with which we can obtain a light neutrino mass spectrum and mixing that are consistent with neutrino oscillation data.

In this model, one of the color-triplet Higgs multiplets lies around an intermediate scale, since its mass is generated by a higher-dimensional operator. As we have seen in section 4.2, this mass, M_{H_C} , is stringently constrained by the current limit on proton decay lifetime provided by the Super-Kamiokande experiments; $M_{H_C} \gtrsim 5 \times 10^{11} \text{ GeV}$ for $\tan \beta \gtrsim 3$. In particular, it is found that the $p \rightarrow \pi^0 \mu^+$ and $p \rightarrow K^0 \mu^+$ modes are most sensitive to the light color-triplet Higgs exchange process. Future proton decay experiments at Hyper-Kamiokande can probe $M_{H_C} \lesssim 10^{12} \text{ GeV}$ in these decay channels. Since these decay channels tend to be subdominant in other GUT models, they offer a promising way to distinguish our model from, e.g., the standard/flipped $SU(5)$ GUT models.

There are several interesting possibilities to extend our model to address other problems. For instance, we may consider a flavor-dependent $U(1)_R$ charge assignment to explain the structure of the Yukawa couplings. In fact, in our model, the ratio $\langle S \rangle / \Lambda$ is predicted to be $\mathcal{O}(0.1)$ for $\Lambda \simeq M_P$, which allows us to take $\langle S \rangle / \Lambda \simeq 0.22 \simeq |V_{us}|$, as often done in models that use the Froggatt-Nielsen mechanism [45] to explain the Yukawa structure. Another direction is to embed our model into a theory based on a larger simple gauge group — e.g., $SO(10)$ — to achieve a complete unification of gauge interactions. These more ambitious model buildings will be explored in future work [46].

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A Model with a discrete R -symmetry

Instead of the $U(1)_R$ symmetry discussed in the text, we may consider a discrete R symmetry to eliminate unwanted terms. In this appendix, we show an example for such models.

We assume that the model considered here possesses a $\mathbb{Z}_{17}$ R symmetry. The matter content of this model is the same as that of the $U(1)_R$ version, as summarized in table 2 with the charge assignment. We also assume this model to have an exact R -parity to suppress dangerous terms such as $F_i \bar{H}$. The superpotential terms allowed by the symmetries in this case are

$$\begin{aligned}
 W = & \frac{1}{4} \lambda_1^{ij} \epsilon_{\alpha\beta\gamma\delta\epsilon} F_i^{\alpha\beta} F_j^{\gamma\delta} h^\epsilon + \sqrt{2} \lambda_2^{ij} F_i^{\alpha\beta} \bar{f}_{j\alpha} \bar{h}_\beta + \lambda_3^{ij} \bar{f}_{i\alpha} \ell_j^c h^\alpha \\
 & + \frac{\lambda_4}{4\Lambda_{\text{DT}}^7} \epsilon_{\alpha\beta\gamma\delta\epsilon} S^7 H^{\alpha\beta} H^{\gamma\delta} h^\epsilon + \frac{1}{4} \lambda_5 \epsilon^{\alpha\beta\gamma\delta\epsilon} \bar{H}_{\alpha\beta} \bar{H}_{\gamma\delta} \bar{h}_\epsilon
 \end{aligned}$$

Fields	Components	SU(5)	U(1)	$\mathbb{Z}_{17R}$	R -parity
F_i	d_i^c, Q_i, ν_i^c	10	+1	1	−
$\bar{f}_i$	u_i^c, L_i	$\bar{5}$	−3	1	−
ℓ_i^c	e_i^c	1	+5	1	−
H	d_H^c, Q_H, ν_H^c	10	+1	8	+
$\bar{H}$	$d_{\bar{H}}^c, Q_{\bar{H}}, \nu_{\bar{H}}^c$	$\bar{10}$	−1	1	+
h	D, H_d	5	−2	0	+
$\bar{h}$	$\bar{D}, H_u$	$\bar{5}$	+2	0	+
S	S	1	0	−2	+

Table 2. The field content and the charge assignments in the $\mathbb{Z}_{17} R$ symmetric model.

$$\begin{aligned}
& + \frac{c_{ij}}{2\Lambda_N^2} S(F_i^{\alpha\beta} \bar{H}_{\alpha\beta})(F_j^{\gamma\delta} \bar{H}_{\gamma\delta}) + \frac{\lambda_H}{4\Lambda_{HS}^5} (H^{\alpha\beta} \bar{H}_{\alpha\beta})^4 + \frac{\lambda_{HS}}{12\Lambda_{HS}^7} (H^{\alpha\beta} \bar{H}_{\alpha\beta})^3 S^4 \\
& + \frac{\lambda'_{HS}}{16\Lambda_{HS}^9} (H^{\alpha\beta} \bar{H}_{\alpha\beta})^2 S^8 + \frac{\lambda''_{HS}}{12\Lambda_{HS}^{11}} (H^{\alpha\beta} \bar{H}_{\alpha\beta}) S^{12} + \frac{\lambda_S}{16\Lambda_{HS}^{13}} S^{16}. \quad (\text{A.1})
\end{aligned}$$

A characteristic feature of this type of models, compared with the $U(1)_R$ models, is that R -axion does not appear after the R symmetry is spontaneously broken. In addition, generically speaking, there are more allowed terms than in the $U(1)_R$ models, which may be advantageous for some cases. For example, we have

$$W_\mu = \frac{\lambda_\mu}{16\Lambda_\mu^{15}} S^{16} h \bar{h}, \quad (\text{A.2})$$

which can generate a μ -term of $\mathcal{O}(m_{\text{SUSY}})$.

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