

Neutrino masses in supersymmetric $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)'$ modelsJunhai Kang,¹ Paul Langacker,¹ and Tianjun Li^{1,2}¹*Department of Physics and Astronomy, University of Pennsylvania, Philadelphia, Pennsylvania 19104-6396, USA*²*School of Natural Science, Institute for Advanced Study, Einstein Drive, Princeton, New Jersey 08540, USA*

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We consider various possibilities for generating neutrino masses in supersymmetric models with an additional $U(1)'$ gauge symmetry. One class of models involves two extra $U(1)' \times U(1)''$ gauge symmetries, with $U(1)''$ breaking at an intermediate scale and yielding small Dirac masses through high-dimensional operators. The right-handed neutrinos N_i^c can naturally decouple from the low energy $U(1)'$, avoiding cosmological constraints. A variant version can generate large Majorana masses for N_i^c and an ordinary seesaw. We secondly consider models with a pair of heavy triplets which couple to left-handed neutrinos. After integrating out the heavy triplets, a small neutrino Majorana mass matrix can be generated by the induced nonrenormalizable terms. We also study models involving the double-seesaw mechanism, in which heavy Majorana masses for N_i^c are associated with the TeV-scale of $U(1)'$ breaking. We give the conditions to avoid runaway directions in such models and discuss simple patterns for neutrino masses.

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I. INTRODUCTION

The possibility of an extra $U(1)'$ gauge symmetry is well-motivated in superstring constructions [1], grand unified theories [2], models of dynamical symmetry breaking [3], little Higgs models [4], and large extra dimensions [5]. In supersymmetric models, an extra $U(1)'$ can provide an elegant solution to the μ problem [6,7], with an effective μ parameter generated by the vacuum expectation value (VEV) of the standard model (SM) singlet field S which breaks the $U(1)'$ symmetry. This is somewhat similar to the effective μ parameter in the next to minimal supersymmetric standard model (NMSSM) [8]. However, with a $U(1)'$ the extra discrete symmetries and their associated cosmological domain wall problems [9] associated with the NMSSM are absent¹. A closely related feature is that the minimal supersymmetric standard model (MSSM) upper bound of M_Z on the tree-level mass of the corresponding lightest MSSM Higgs scalar is relaxed, both in models with a $U(1)'$ and in the NMSSM, because of the Yukawa term hSH_1H_2 in the superpotential [8] and the $U(1)'$ D -term [11]. More generally, for specific $U(1)'$ charge assignments for the ordinary and exotic fields one can simultaneously ensure the absence of anomalies; that all fields of the TeV-scale effective theory are chiral, avoiding a generalized μ problem; and the absence of dimension-4 proton decay operators [12]. $U(1)'$ models can also be consistent with gauge unification and may have implications in electroweak baryogenesis [13], cold dark matter [14], rare B decays [15], and nonstandard Higgs potentials [16].

There are stringent limits from direct searches at the Tevatron [17] and from indirect precision tests at the Z -pole, at LEP 2, and from weak neutral current experiments [18]. The constraints depend on the particular Z'

couplings, but in typical models one requires $M_{Z'} > (500\text{--}800)$ GeV and the $Z - Z'$ mixing angle $\alpha_{Z-Z'}$ to be smaller than a few $\times 10^{-3}$. Thus, explaining the $Z - Z'$ mass hierarchy is important. Recently, we proposed a supersymmetric model with a string-motivated secluded $U(1)'$ -breaking sector, where the squark and slepton spectra can mimic those of the MSSM, the electroweak symmetry breaking is driven by relatively large A terms, and a large Z' mass can be generated by the VEVs of additional SM singlet fields that are charged under the $U(1)'$ [19].

On the other hand, very light left-handed neutrinos, which has been confirmed from the solar and atmospheric neutrino experiments and KamLAND experiment, is a mystery in nature. Possible scenarios [20] with tiny left-handed neutrino masses include extensions of the SM at a low energy scale, for example, the Zee model [21], in which the left-handed neutrino masses are generated at loop level; supersymmetric models with lepton number and R-parity violation [22], which can include both tree and loop effects; double or extended (i.e., TeV-scale) seesaw models [23]; or models including large extra dimensions [24]. Mechanisms involving high energy scales include the canonical seesaw mechanism, in which heavy right-handed Majorana neutrinos have masses of the order $10^{10}\text{--}10^{16}$ GeV [25]; models involving heavy Higgs triplets [26–28]; and models in which small neutrino Dirac masses are generated by high-dimensional operators [29,30].

In this paper, we consider the possibilities for small neutrino masses in supersymmetric $U(1)'$ models. The $U(1)'$ symmetry affects some of the above mechanisms which can generate the tiny neutrino masses. In particular, right-handed neutrinos could not acquire Majorana masses at a scale much larger than the $U(1)'$ -breaking scale unless they are not charged under $U(1)'$, thus forbidding a canonical high-scale seesaw mechanism in many TeV-scale $U(1)'$ models. Another implication involves the right-handed

¹For other solutions, see [10].

neutrinos in models with small neutrino Dirac masses. In the SM these are harmless cosmologically because they are essentially sterile (except for negligible Higgs couplings and mass effects) and are not produced in significant numbers prior to big bang nucleosynthesis (BBN). However, in $U(1)'$ models the right-handed neutrinos can be produced by these Z' interactions [unless their $U(1)'$ charge vanishes], leading to stringent constraints on the Z' mass [31]. Other, comparable, constraints follow from supernova cooling [32].

We discuss a number of possibilities for neutrino masses in $U(1)'$ models. In Sec. II, we consider the possibility of small Dirac masses. We assume that elementary renormalizable neutrino Yukawa couplings are forbidden by the extra gauge symmetry, other symmetries, or string selection rules, but that effective neutrino Yukawa couplings are generated by nonrenormalizable terms after certain SM singlet fields acquire intermediate-scale VEVs. Essentially speaking, this is a generalization of the Froggatt-Nielsen model [33]. We consider models with two additional $U(1)' \times U(1)''$ gauge symmetries, with $U(1)''$ breaking at an intermediate scale and $U(1)'$ at the TeV-scale. The intermediate-scale $U(1)''$ and the associated high-dimensional operators can account for small neutrino Dirac masses. It can occur naturally that after the intermediate-scale $U(1)''$ breaking, the right-handed neutrinos are neutral under the TeV-scale $U(1)'$ so as to avoid the BBN and supernova constraints. The existence of two extra $U(1)$ s is partly motivated by E_6 grand unification, since E_6 can be broken down to the SM gauge group with two additional $U(1)$ s. However, we only use the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)' \times U(1)'' \subset E_6$ quantum number assignments for the ordinary and exotic particles in **27** and **27*** of E_6 to construct an example of an anomaly-free model, and we do not consider the full E_6 model². We describe how the symmetry breaking pattern can be realized assuming the Higgs fields from **27** and **27*** of E_6 . We also give an example of how the small neutrino Dirac masses can be generated in a model with a TeV-scale secluded $U(1)'$ -breaking sector as in [19]. In that case, however, the decoupling of the right-handed neutrinos requires the introduction of singlets not belonging to simple E_6 representations. In these models there are no allowed couplings that can generate large Majorana masses for the right-handed neutrinos at the intermediate scale. However, we also consider a variant case in which there are allowed couplings which can generate large effective Majorana masses for the right-handed neutrinos through the intermediate-scale $U(1)''$ breaking, leading to a traditional seesaw.

Instead of generating small neutrino Dirac masses from Yukawa couplings with doublet Higgs fields and high-

²The Yukawa relations for a full E_6 theory would lead to rapid proton decay for a low $U(1)'$ breaking scale.

dimensional operators, one can generate small neutrino Majorana masses through their couplings with triplet fields. We propose two models involving a pair of heavy triplets. The mass for the triplets is about 10^{14} GeV for the first model, and about 10^8 GeV for the second. After the electroweak symmetry breaking, the triplets obtain very small VEVs and give a realistic left-handed neutrino Majorana mass matrix. Equivalently, one can integrate out the heavy triplets and obtain a low energy neutral Higgs potential that is the same as that in our previous model [19], up to negligible corrections, with the left-handed neutrino Majorana mass matrix generated by the induced nonrenormalizable terms.

Yet another possibility is the double-seesaw mechanism. If the right-handed neutrinos are charged under $U(1)'$, they may acquire Majorana masses at the $U(1)'$ -breaking scale. We consider a model with the double-seesaw mechanism, in which the neutrino masses are suppressed by two powers of the TeV-scale masses. The neutrino Yukawa couplings can be of order 10^{-3} , i.e., the neutrino Dirac masses are of the order of the muon mass. The double-seesaw mechanism has been discussed previously for one family in Ref. [23]; here, we generalize it to three families in the context of $U(1)'$ models.

We slightly modify the model in Ref. [19] by introducing three right-handed neutrinos and three SM singlets. Runaway directions can be avoided by imposing suitable conditions on the soft terms. The vacuum is the same as in [19], so the previous discussions on the $Z - Z'$ mass hierarchy and the particle spectrum still hold. The active neutrino mass matrix is $M_D(M_V^{-1})^T M_B M_V^{-1} M_D^T$, where M_D is the 3×3 Dirac mass matrix, and M_V and M_B are 3×3 matrices defined in Sec. IV. Because the typical mass scale for M_V is TeV, the active neutrinos may have realistic masses and mixings if the typical mass scales for M_B and M_D are about 0.1 GeV. We show that normal, inverted and degenerate textures can be achieved from reasonable assumptions about M_D , M_V , and M_B .

This paper is organized as follows: in Sec. II we consider the supersymmetric $U(1)' \times U(1)''$ models with $U(1)''$ breaking at the intermediate scale to generate small neutrino Dirac masses. We discuss two models with a pair of heavy triplets in Sec. III. In Sec. IV, we consider the supersymmetric $U(1)'$ model with the double-seesaw mechanism. Our discussions and conclusions are given in Sec. V. We discuss the runaway directions for the double-seesaw model in the appendix.

II. GENERATING NEUTRINO MASSES FROM INTERMEDIATE-SCALE $U(1)''$ BREAKING

We first consider the possibility of generating small neutrino Dirac masses in $U(1)' \times U(1)''$ models where the $U(1)'$ and $U(1)''$ are broken at the TeV-scale and intermediate scale, respectively. In this case, we must consider the constraints from BBN [31]. If the right-

handed neutrinos are charged under the $U(1)'$, they will couple to other particles through the exchange of $U(1)'$ gauge boson. They must decouple well before the BBN epoch so as to avoid the BBN constraints from the predicted ${}^4\text{He}$ abundance. Either the $U(1)'$ must be broken at a high scale, typically above 5 TeV, or the right-handed neutrinos are neutral under the $U(1)'$. Complementary constraints follow from supernova cooling [32]. Here we show that the N_i^c decouplings can occur naturally in certain $U(1)' \times U(1)''$ models [34].

We consider a model with the gauge group $G_{SM} \times U(1)' \times U(1)''$, where G_{SM} is the SM gauge group. The $U(1)''$ is broken at an intermediate scale around 10^{10} GeV, the right-handed neutrinos are left neutral under the $U(1)'$, the small neutrino Dirac masses are due to high-dimensional operators associated with the intermediate scale, and the $U(1)' \times U(1)''$ symmetry forbids both elementary Majorana masses for the right-handed neutrinos and also renormalizable-level interactions that could generate their large effective Majorana masses at the intermediate scale.

The $U(1)''$ can be broken at the intermediate scale if it is associated with a potential which is F- and D- flat at tree level. For example, if we introduce one pair of vectorlike SM singlets S_1 and S_1^* , the F-flatness implies a tree-level potential

$$V(S_1, S_1^*) = m_{S_1}^2 |S_1|^2 + m_{S_1^*}^2 |S_1^*|^2 + \frac{g'^2 Q''^2}{2} (|S_1|^2 - |S_1^*|^2)^2, \quad (1)$$

where g' is the $U(1)''$ gauge coupling constant and Q'' is the $U(1)''$ charge for S_1 .

For $m_{S_1}^2 + m_{S_1^*}^2 < 0$, there will be a runaway direction along the D-flat direction $|\langle S_1 \rangle| = |\langle S_1^* \rangle|$. However, the potential will be stabilized by loop corrections or high-dimensional operators, so that the S_1 and S_1^* will obtain intermediate-scale VEVs. Neutrino Dirac masses could be generated by high-dimensional operators, such as

$$W \sim H_2 L_i N_j^c \left(\frac{S}{M_{Pl}} \right)^{P_D}, \quad (2)$$

where L_i and N_j^c are the superfields, respectively, corresponding to the lepton doublets and right-handed neutrinos, and $M_{Pl} \sim 10^{19}$ GeV is the Planck scale. The S field can be S_1 or S_1^* or any combinations that are allowed by gauge invariance and other symmetries of the four-dimensional theory, and by string selection rules. It is reasonable that in some models neutrino mass terms may occur in higher order than those for the quarks and charged leptons, leading to naturally small neutrino Dirac masses. Choosing proper S field VEVs and powers P_D , one can obtain a realistic neutrino mass spectrum. However, without a more detailed construction, there is no predictive

power for the type of neutrino hierarchy and the mixing angles.

As an example, we consider how this mechanism can be realized using the $U(1)' \times U(1)''$ charges associated with the **27** representation of the E_6 gauge group. We show that for the appropriate signs of soft-supersymmetry breaking parameters, small neutrino Dirac masses can be generated, with the N_i^c naturally decoupling from the TeV-scale $U(1)'$, satisfying the BBN and supernova constraints. We also consider how to incorporate small neutrino Dirac masses in the secluded $U(1)'$ model. In this case, the decoupling of the N_i^c from the TeV-scale $U(1)'$ does not occur if all the particles arise from the simple E_6 representations, but could in a more general context. We also study the possibility of generating large Majorana masses for right-handed neutrinos in the context of intermediate-scale $U(1)''$ breaking.

A. Neutrino masses in models with E_6 particle content

Models with the gauge group $G_{SM} \times U(1)' \times U(1)''$ may appear in grand unification theory with the E_6 gauge group, since E_6 can be broken down to the SM through

$$E_6 \rightarrow SO(10) \times U(1)_\psi \rightarrow SU(5) \times U(1)_\chi \times U(1)_\psi. \quad (3)$$

The $U(1)_\chi$ and $U(1)_\psi$ charges for the particles in the **27** representation are given in Table I. The representations of E_6 are automatically anomaly free, so it is an example of a consistent model with additional $U(1)$ s. In a full E_6 grand unified theory, the two extra $U(1)$ s would have to be broken at the GUT scale, because otherwise they would prevent the exotic D -quark partners and the Higgs doublets in the **27**s from acquiring large masses, and the D -scalars could mediate rapid proton decay. Nevertheless, it is convenient to use the $G_{SM} \times U(1)_\chi \times U(1)_\psi \subset E_6$ quantum number assignments for the ordinary and exotic particles in **27** and **27*** to construct an example of an anomaly-free $U(1)' \times$

TABLE I. Decomposition of the E_6 fundamental representation **27** under $SO(10)$, $SU(5)$, and $U(1)$ s for the particles in the **27**. $U(1)_\chi$ and $U(1)_\psi$ are orthogonal to each other with $Q_\chi = 0$ for S_L . Q_1 and Q_2 are, respectively, the particle charges under the $U(1)_1$ and $U(1)_2$ which are orthogonal with $Q_1 = 0$ for \bar{N} . Q is the particle charge under the $U(1)'$ in an anomaly-free supersymmetric $U(1)'$ model with a secluded $U(1)'$ -breaking sector [13].

$SO(10)$	$SU(5)$	$2\sqrt{10}Q_\chi$	$2\sqrt{6}Q_\psi$	$2\sqrt{15}Q$	$2\sqrt{6}Q_2$	$2\sqrt{10}Q_1$
16	$10(u, d, \bar{u}, \bar{e})$	-1	1	-1/2	1	1
	$\bar{5}(\bar{d}, \nu, e)$	3	1	4	-2	2
	$1\bar{N}$	-5	1	-5	4	0
10	$5(D, H'_u)$	2	-2	1	-2	-2
	$\bar{5}(\bar{D}, H'_d)$	-2	-2	-7/2	1	-3
1	$1S_L$	0	4	5/2	1	5

U(1)'' model, even though the rest of E_6 structure, such as the Yukawa relations, is violated. This is typical in string constructions [35], which often do not respect the E_6 Yukawa relations (that would be responsible for proton decay), or which may lead to more complicated U(1)' charge assignments and exotic structure.

We first give an example of a U(1)' \times U(1)'' model in which U(1)'' is broken at the intermediate scale while the U(1)', which is broken at the TeV-scale, decouples from the right-handed neutrinos. There are two SM singlets in the **27**, \tilde{N} , which we will identify as the right-handed neutrinos N_i^c , and S_L . There is only one linear combination of U(1) $_{\chi}$ and U(1) $_{\psi}$ in which the N_i^c fields are neutral, that is the U(1) $_1$ shown in Table I³. U(1) $_2$ is the other linear combination, which is orthogonal to U(1) $_1$. Here, to avoid confusion, we consider the U(1) $_1$ and U(1) $_2$ as the U(1)' and U(1)'', respectively. We can naturally break the U(1) $_2$ at a high scale by giving large VEVs to the scalar components ($\tilde{\nu}_R^*$ and $\tilde{\nu}_R$) of a pair of the vectorlike superfields ν_R^* and ν_R whose quantum numbers are the same as those of the N_i^c and its Hermitian conjugate⁴. The F- and D-flatness can be preserved and the U(1) $_1$ unbroken until the TeV-scale. To have D-flat directions, we introduce two pairs of the vectorlike fields (ν_R, ν_R^*) and (S_L, S_L^*) from the singlets of (**27**, **27**^{*}) in addition to the SM fermions from three complete **27**-plets. Then, the D-term potential is

$$V_{\chi} + V_{\psi} = \frac{g'^2}{2} \left[\frac{5}{2\sqrt{10}} (|\tilde{\nu}_R|^2 - |\tilde{\nu}_R^*|^2) \right]^2 \quad (4)$$

$$+ \frac{g'^2}{2} \left[\frac{1}{\sqrt{24}} (-|\tilde{\nu}_R|^2 + |\tilde{\nu}_R^*|^2 - 4|S_L|^2 + 4|S_L^*|^2) \right]^2, \quad (5)$$

where a sum over each type of scalar is implied, and we have assumed equal gauge couplings for simplicity. The potential is clearly D-flat for $|\langle \tilde{\nu}_R \rangle|^2 = |\langle \tilde{\nu}_R^* \rangle|^2 \equiv |\langle \tilde{\nu} \rangle|^2$ and $|\langle S_L \rangle|^2 = |\langle S_L^* \rangle|^2 \equiv |\langle S \rangle|^2$. We assume that the potential is also F-flat along this direction. The potential along the flat direction is then

$$V = m_{\tilde{\nu}}^2 |\tilde{\nu}^2| + m_S^2 |S|^2, \quad (6)$$

where $m_{\tilde{\nu}}^2$ and m_S^2 are, respectively, the sum of the mass-squares of the $\tilde{\nu}_R$ and $\tilde{\nu}_R^*$, and that of the S_L and S_L^* , which we assume are typical soft-supersymmetry breaking scale. For $m_S^2 > 0$ and $m_{\tilde{\nu}}^2 < 0$, the breaking will occur along the

³The U(1) $_1$ charge assignment has been found in Ref. [36] from different motivations, i.e., to explain the tiny active neutrino masses via the high-scale seesaw mechanism or allow for the possibility of leptogenesis. However, leptogenesis is not required since electroweak baryogenesis in U(1)' models is a much more viable possibility than in the MSSM [13].

⁴The scalar components \tilde{N}_i^c of the right-handed neutrino superfields N_i^c should not acquire large VEVs to avoid large lepton-Higgsino mixings. We therefore assume that the scalars \tilde{N}_i^c have positive soft mass-squares, while the $\tilde{\nu}_R$ and $\tilde{\nu}_R^*$ can acquire large VEVs.

D-flat direction for $|\langle \tilde{\nu}_R \rangle| = |\langle \tilde{\nu}_R^* \rangle|$ very large, with the potential ultimately stabilized by loop corrections or high-dimensional operators [29]. However, since $m_S^2 > 0$, S_L and S_L^* will acquire (usually different) TeV-scale VEVs not associated with the flat direction.

For arbitrary VEVs, the mass terms for extra gauge bosons are

$$L = g'^2 \left(-\frac{5}{2\sqrt{10}} Z_{\chi} + \frac{1}{\sqrt{24}} Z_{\psi} \right)^2 (|\tilde{\nu}_R|^2 + |\tilde{\nu}_R^*|^2) + g'^2 \left(\frac{4}{\sqrt{24}} Z_{\chi} \right)^2 (|S_L|^2 + |S_L^*|^2). \quad (7)$$

For the breaking pattern described above, this will imply that $\sqrt{2/3} Z_2 = -[5/(2\sqrt{10})] Z_{\chi} + (1/\sqrt{24}) Z_{\psi}$ will acquire a superheavy mass, while the orthogonal combination $\sqrt{2/3} Z_1 = (1/\sqrt{24}) Z_{\chi} + [5/(2\sqrt{10})] Z_{\psi}$ will remain at the TeV-scale. Z_1 decouples from N_i^c and therefore evades the nucleosynthesis and supernova constraints.

As an example of a high-dimensional operator to stabilize the potential in Eq. (6), let us consider

$$W \sim c \frac{(\nu_R \nu_R^*)^2}{M_{Pl}}. \quad (8)$$

The ν_R and ν_R^* fields will obtain VEVs around $10^{10}/\sqrt{c}$ to $10^{11}/\sqrt{c}$ GeV in this case⁵. The small neutrino Dirac mass terms may be generated through

$$W \sim H_2 L_i N_j^c \frac{\nu_R \nu_R^*}{M_{Pl}^2}, \quad (9)$$

which is typically of order $10^{-6}/c$ to $10^{-5}/c$ eV. A small $c \sim 10^{-3}-10^{-4}$ would yield appropriate neutrino masses. Such a value for c could be generated if the operator in Eq. (8) was itself due to a high-dimensional operator involving additional fields with VEVs close to M_{Pl} , e.g., associated with an anomalous U(1)' [37].

B. Small neutrino Dirac masses in a U(1)' model with a secluded sector

The TeV-scale U(1)' could appear in a model with a secluded U(1)'-breaking sector as in [19]. This model can solve the supersymmetric μ problem, contribute to electroweak baryogenesis, and yield a $Z - Z'$ hierarchy and small mixing angle. In this subsection, we show how to extend the U(1)' \times U(1)'' model discussed above to incorporate a secluded sector.

The superpotential for the Higgs sector in the secluded U(1)' model is

$$W_H = h S H_1 H_2 + \lambda S_1 S_2 S_3, \quad (10)$$

⁵Alternatively, such an intermediate scale could be generated by loop corrections to the effective potential, which would render the running $m_{\tilde{\nu}}^2$ positive at the intermediate scale.

where the Yukawa coupling h is associated with the effective μ term and the potential has a runaway direction for $\lambda \rightarrow 0$; the S and S_i fields are SM singlets, with $U(1)'$ charge assignments

$$Q_S = -Q_{S_1} = -Q_{S_2} = \frac{1}{2}Q_{S_3}, \quad Q_{H_1} + Q_{H_2} + Q_S = 0. \quad (11)$$

For a sufficiently small value of λ , the Z' mass can be arbitrarily large. For example, if $h \sim 10\lambda$, one can generate a $Z - Z'$ mass hierarchy in which the Z' mass is of order one TeV [19].

A $U(1)'$ model with a secluded sector using the E_6 particle contents and charge assignments was constructed in [13]. In that model, it was assumed that the four SM singlets S, S_1, S_2, S_3 are the S_L, S_L^*, S_L^* and \tilde{N}^* , respectively, in two pairs of **27** and **27***. Choosing a special combination of $U(1)_\chi$ and $U(1)_\psi$ (see the Q charge assignments in Table I), the charge relations in Eq. (11) are satisfied. However, in that model, the right-handed neutrinos N_i^c are charged under the low energy $U(1)'$, and will be constrained by the BBN and supernova data if the neutrinos have small Dirac masses.

We can instead consider charge assignments such that the right-handed neutrinos N_i^c will be neutral under the TeV-scale $U(1)'$, i.e., the charge assignments of Q_1 and Q_2 in Table I. As discussed in the last subsection, the scalar components of the superfields ν_R^* and ν_R with the same quantum numbers as those of N_i^c and its Hermitian conjugate will obtain intermediate-scale VEVs and break the $U(1)_2$ and leave a TeV-scale $U(1)_1$. To incorporate the secluded $U(1)'$ model, one must introduce SM singlets that satisfy the $U(1)'$ charge relations in Eq. (11). This is not possible for S or S_i fields belonging to the (**27**, **27***) or other low-dimensional representations of E_6 . However, recalling that we are using E_6 only as an example of an anomaly-free construction, it is not unreasonable to consider the possibility of charge assignments for SM singlets that do not correspond to E_6 , as long as they are vectorlike pairs so as to avoid anomalies. For example, we can assume the SM singlets S, S_1, S_2 are the S_L, S_L^*, S_L^* , respectively, in two pairs of **27** and **27***, and also introduce one pair of vectorlike fields S_3 and S_3^* with $U(1)'$ charge $Q_{S_3} = -Q_{S_3^*} = 2Q_{S_L}$. In this way, we can generate small neutrino Dirac masses from the intermediate-scale $U(1)''$ breaking in the secluded $U(1)'$ -breaking model.

C. Large majorana masses for right-handed neutrinos

The above discussions in the $U(1)' \times U(1)''$ models concentrated on generating small neutrino Dirac masses from the intermediate-scale $U(1)''$ breaking. One can also generate the large Majorana masses for the right-handed neutrinos, to yield the ordinary seesaw mechanism.

Let us consider three right-handed neutrinos N_i^c and one pair of vectorlike fields S and S^* , with charges,

$$Q'_{N_i^c} = Q'_S = 0, \quad Q''_{N_i^c} = -\frac{1}{2}Q''_S, \quad (12)$$

where Q' and Q'' are the particle charges under the TeV-scale $U(1)'$ and intermediate-scale $U(1)''$, respectively. Then, the superpotential is

$$W \sim \frac{1}{M_{Pl}^{2k-3}}(SS^*)^k + SN_i^c N_j^c. \quad (13)$$

We also introduce the soft-supersymmetry breaking terms

$$V \sim m_{N_i^c}^2 |\tilde{N}_i^c|^2 + m_S^2 |S|^2 + m_{S^*}^2 |S^*|^2, \quad (14)$$

where \tilde{N}_i^c is the scalar component of the superfield N_i^c . If we assume $m_{N_i^c}^2 > 0$ while $m_S^2 + m_{S^*}^2 < 0$, the VEVs of the S and S^* fields will be driven to nonzero values, while those of the \tilde{N}_i^c fields will be zero. The D-flat direction will ensure $\langle S \rangle = \langle S^* \rangle$. The potential will be stabilized by the high-dimensional operators, which determines $\langle S \rangle \sim (m_S M_{Pl}^{2k-3})^{1/(2k-2)}$. Taking, for example, $k = 3$ will yield $\langle S \rangle \sim 10^{14}$ GeV. These VEVs will give Majorana masses to the N_i^c fields of the same order, allowing an ordinary seesaw mechanism.

III. HIGGS TRIPLET MODELS

A number of authors have considered models in which small neutrino Majorana masses can be generated by coupling two lepton doublets to an $SU(2)_L$ -triplet T with weak hypercharge $Y = 1$. Early versions of the triplet models [26] assumed spontaneous lepton number violation. These are excluded by the invisible Z width, which would be increased equivalent to two extra neutrino species by Z decaying into the Goldstone boson (Majoron) and a light scalar. However, more recent scenarios [27] avoid this difficulty by coupling T to the Higgs doublets as well, which breaks lepton number explicitly. These couplings ensure that T^0 acquires a tiny VEV if T is given a very large mass, or equivalently lead to the suppressed high-dimensional operators if T is integrated out. Such models are sometimes referred as the type II seesaw mechanism. Supersymmetric versions have been constructed [28], and there are special constraints when this mechanism is embedded in string constructions, as discussed in [38]. Here, we show that the type II seesaw mechanism can be applied in the supersymmetric $U(1)'$ models.

We consider the supersymmetric $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)'$ models with a pair of very heavy triplets T and \tilde{T} , which can have very small VEVs for the charge zero components (T^0 and \tilde{T}^0) after electroweak symmetry breaking, and give the needed neutrino Majorana masses and mixings. The quantum numbers for T and \tilde{T} are $(1, 3, 1)$ and $(1, 3, -1)$ under the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetry. To be concrete, we integrate out the heavy triplets T and \tilde{T} , and find that the low energy neutral Higgs potential is almost the same as that in

Ref. [19]. Moreover, with suitable Yukawa couplings for the lepton doublets and triplets, a realistic neutrino Majorana mass matrix can be generated by nonrenormalizable terms. For simplicity, we only consider the neutral Higgs potential and the Yukawa couplings for the left-handed neutrino Majorana masses.

A. Model I

In model I, the $U(1)'$ charges for the Higgs fields, triplets, and lepton doublets are

$$Q_{H_1} + Q_{H_2} + Q_S = 0, \quad Q_L \equiv Q_{L_i} = -Q_{H_2}, \quad (15)$$

$$Q_T = -Q_{\bar{T}} = 2Q_{H_2}. \quad (16)$$

We choose the superpotential

$$W = hSH_1H_2 + \lambda_u H_2 \bar{T} H_2 + y_{ij} L_i T L_j + mT\bar{T}, \quad (17)$$

where m is the mass for T and \bar{T} , which is about 10^{14} GeV. We do not need to introduce right-handed neutrinos. Even if there are right-handed neutrinos, the Yukawa terms $L_i H_2 N_j^c$ are assumed to be forbidden by the $U(1)'$, or the other symmetries or the underlying string constructions. The F -term neutral scalar potential is

$$V_F = h^2 |H_1^0 H_2^0|^2 + h^2 |SH_2^0|^2 + |hSH_1^0 + 2\lambda_u \bar{T}^0 H_2^0|^2 + |\lambda_u H_2^0 H_2^0 + mT^0|^2 + m^2 |\bar{T}^0|^2, \quad (18)$$

and the D -term potential is

$$V_D = \frac{G^2}{8} (|H_2^0|^2 - |H_1^0|^2 + 2|T^0|^2 - 2|\bar{T}^0|^2)^2 + \frac{1}{2} g_{Z'}^2 (Q_S |S|^2 + Q_{H_1} |H_1^0|^2 + Q_{H_2} |H_2^0|^2 + 2Q_{H_2} |T^0|^2 - 2Q_{H_2} |\bar{T}^0|^2)^2, \quad (19)$$

where $G^2 = g_1^2 + g_2^2$; g_1 , g_2 , and $g_{Z'}$ are the coupling constants for $U(1)_Y$, $SU(2)_L$ and $U(1)'$; and Q_ϕ is the $U(1)'$ charge of the field ϕ .

We also consider the Yukawa coupling for the neutrinos

$$\mathcal{L}_{\text{Yukawa}} = -\frac{1}{2} y'_{ij} \nu_i T^0 \nu_j + \text{H.c.}, \quad (20)$$

where ν_i are the left-handed neutrinos, and $y'_{ij} = y_{ij}(1 + \delta_{ij})$ in which δ_{ij} is equal to 1 or 0 for $i = j$ or $i \neq j$, respectively.

After electroweak symmetry breaking, i.e., $H_1^0 \neq 0$ and $H_2^0 \neq 0$, the F -terms for H_2^0 , T^0 , and \bar{T}^0 cannot be zero simultaneously. The T^0 and \bar{T}^0 will acquire very small VEVs

$$\langle T^0 \rangle \simeq -\frac{\lambda_u}{m} \langle H_2^0 \rangle \langle H_2^0 \rangle, \quad \langle \bar{T}^0 \rangle \simeq -\frac{2h\lambda_u^*}{m^2} \langle S \rangle \langle H_1^0 \rangle \langle H_2^{0*} \rangle. \quad (21)$$

There are no experimental constraints on the VEVs of T^0 and \bar{T}^0 in this range (i.e., much smaller than the electro-

weak scale). The left-handed neutrino Majorana mass terms are given by

$$\mathcal{L}_{\text{Yukawa}} = \frac{\lambda_u}{2m} y'_{ij} \langle H_2^0 \rangle \langle H_2^0 \rangle \nu_i \nu_j + \text{H.c.} \quad (22)$$

Alternatively, we can integrate out the T^0 and \bar{T}^0 because they are heavy. Their equations of motion are

$$m\lambda_u H_2^0 H_2^0 + \frac{1}{2} y'_{ij} \nu_i \nu_j + (m^2 + 4\Delta_{EW} + 4\Delta_{Z'} Q_{H_2}) T^0 + (G^2 + 4g_{Z'}^2 Q_{H_2}^2) |T^0|^2 T^0 = 0, \quad (23)$$

$$2\lambda_u^* h S H_1^0 H_2^{0*} + (m^2 + 4\lambda_u^2 |H_2^0|^2 - 4\Delta_{EW} - 4\Delta_{Z'} Q_{H_2}) \bar{T}^0 + (G^2 + 4g_{Z'}^2 Q_{H_2}^2) |\bar{T}^0|^2 \bar{T}^0 = 0, \quad (24)$$

where

$$\Delta_{EW} = \frac{G^2}{8} (|H_2^0|^2 - |H_1^0|^2), \quad (25)$$

$$\Delta_{Z'} = \frac{1}{2} g_{Z'}^2 \left(Q_S |S|^2 + Q_{H_1} |H_1^0|^2 + Q_{H_2} |H_2^0|^2 + \sum_{i=1}^3 Q_{S_i} |S_i|^2 \right). \quad (26)$$

The terms proportional to $|T^0|^3$ and $|\bar{T}^0|^3$ are very small due to the large m . Thus,

$$T^0 \simeq -\frac{m\lambda_u H_2^0 H_2^0 + \frac{1}{2} y'_{ij} \nu_i \nu_j}{m^2 + 4\Delta_{EW} + 4\Delta_{Z'} Q_{H_2}} \sim -\frac{\lambda_u H_2^0 H_2^0}{m}, \quad (27)$$

and the resulting nonrenormalizable neutrino mass terms are

$$\mathcal{L}_{\text{Yukawa}} = \frac{1}{2} \frac{\lambda_u y'_{ij} \nu_i \nu_j H_2^0 H_2^0}{m} + \text{H.c.} \quad (28)$$

The neutrino mass (m_ν) scale is about 0.05 eV, implying $m \sim 10^{14}$ GeV. With suitable Yukawa couplings y_{ij} , one can obtain a realistic left-handed neutrino Majorana mass matrix. Of course, the $U(1)'$ symmetry does not by itself constrain the form of y_{ij} or lead to a prediction for the form of the mass hierarchy and mixings. The low energy neutral Higgs potential is just that in Ref. [19] up to negligible corrections of order $(M_{Z'}/m)^2 \sim 10^{-24}$.

B. Model II

In model II, the $U(1)'$ charges for the Higgs fields, triplets and lepton doublets are

$$Q_{H_1} + Q_{H_2} + Q_S = 0, \quad Q_{L_i} = Q_L = Q_{H_1}, \quad (29)$$

$$Q_T = -Q_{\bar{T}} = -2Q_{H_1}, \quad (30)$$

and the superpotential is

$$W = hSH_1H_2 + \lambda_d H_1TH_1 + y_{ij}L_iTL_j + mT\bar{T}, \quad (31)$$

where m is the mass for T and \bar{T} , around 10^8 GeV.

Similar to the last subsection, after electroweak symmetry breaking, i.e., $H_1^0 \neq 0$ and $H_2^0 \neq 0$, the F -terms for H_1^0 , T^0 and \bar{T}^0 cannot be zero simultaneously, and the T^0 and \bar{T}^0 will have very small VEVs

$$\langle T^0 \rangle \simeq -\frac{2\lambda_d^* h \langle S \rangle \langle H_2^0 \rangle \langle H_1^{0*} \rangle}{m^2}, \quad \langle \bar{T}^0 \rangle \simeq -\frac{\lambda_d}{m} \langle H_1^0 \rangle \langle H_1^0 \rangle. \quad (32)$$

The left-handed neutrino Majorana masses are given by

$$\mathcal{L}_{\text{Yukawa}} = \frac{1}{m^2} y'_{ij} \lambda_d^* h \langle S \rangle \langle H_2^0 \rangle \langle H_1^{0*} \rangle \nu_i \nu_j + \text{H.c.}, \quad (33)$$

which can also be obtained by integrating out T^0 and \bar{T}^0 . $m_\nu \sim 0.05$ eV can be obtained for $m \sim 10^8$ GeV. Suitable Yukawa couplings y_{ij} can yield a realistic left-handed neutrino Majorana mass matrix, with negligible small corrections of order $(M_Z/m)^2 \sim 10^{-12}$ to the low energy Higgs potential.

IV. THE DOUBLE-SEESAW MECHANISM

Another possibility for small neutrino Majorana masses is the double or extended seesaw mechanism [23]. Typically, the large scale in such models is only of the order TeV. However, the light neutrino masses are suppressed by two or more powers of this scale and sometimes small scales in the numerator. Such constructions have been suggested, e.g., in the context of superstring model buildings [39], in which it is difficult to generate a normal seesaw [40]. They are also a viable possibility in the $U(1)'$ models, in which the TeV-scale masses may be associated with the $U(1)'$ breaking scale. In this section, we show that the secluded sector model can be extended to include the double-seesaw mechanism, without introducing unwanted runaway directions, and that one can obtain the normal and inverted hierarchies, and the degenerate scenarios, for neutrino masses.

We consider the supersymmetric $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)'$ model with two Higgs doublets (H_1 and H_2), four Higgs singlets (S , S_1 , S_2 and S_3), and three extra singlets (B_1 , B_2 and B_3). Assuming the $U(1)'$ charges satisfy the equations

$$Q_{N_i^c} \equiv Q_{N^c} = -\frac{3}{2} Q_S, \quad Q_{B_i} \equiv Q_B = -\frac{1}{2} Q_S, \quad (34)$$

$$Q_{L_i} \equiv Q_L = -Q_{H_2} + \frac{3}{2} Q_S, \quad (35)$$

as well as these in Eq. (11), we choose the superpotential

$$W = hSH_1H_2 + \lambda S_1S_2S_3 + d_{ij}SB_iB_j + e_{ij}S_3N_i^cB_j + y_{ij}H_2N_i^cL_j, \quad (36)$$

where h , λ , d_{ij} , e_{ij} and y_{ij} are Yukawa couplings, and we assume that $d_{ii} = 0$, motivated by string constructions. The corresponding F -term scalar potential is

$$V_F = h^2|S|^2|H_2|^2 + |hSH_1 + y_{ij}\tilde{N}_i^c\tilde{L}_j|^2 + |hH_1H_2 + d_{ij}\tilde{B}_i\tilde{B}_j|^2 + \lambda^2(|S_2|^2 + |S_1|^2)|S_3|^2 + |\lambda S_1S_2 + e_{ij}\tilde{N}_i^c\tilde{B}_j|^2 + \sum_{i=1}^3 |e_{ij}S_3\tilde{B}_j + y_{ij}H_2\tilde{L}_j|^2 + \sum_{j=1}^3 |y_{ij}H_2\tilde{N}_i^c|^2 + \sum_{j=1}^3 |d_{ij}S\tilde{B}_i + e_{ij}S_3\tilde{N}_i^c|^2, \quad (37)$$

where for a supermultiplet ϕ which is not a Higgs doublet (H_1 or H_2) or singlet field (S or S_i), we denote its scalar component as $\tilde{\phi}$. The D -term scalar potential for the fields that are $SU(3)$ singlets and neutral under $U(1)_Y$ is

$$V_D = \frac{G^2}{8} \left(|H_2^0|^2 - |H_1^0|^2 - \sum_{i=1}^3 |\tilde{\nu}_i|^2 \right)^2 + \frac{1}{2} g_{Z'}^2 \left[Q_S |S|^2 + Q_{H_1} |H_1^0|^2 + Q_{H_2} |H_2^0|^2 + \sum_{i=1}^3 Q_{S_i} |S_i|^2 + \sum_{i=1}^3 (Q_N |\tilde{N}_i^c|^2 + Q_L |\tilde{\nu}_i|^2 + Q_B |\tilde{B}_i|^2) \right]^2. \quad (38)$$

In addition, we introduce the supersymmetry breaking soft terms

$$V_{\text{soft}} = m_{H_1}^2 |H_1|^2 + m_{H_2}^2 |H_2|^2 + m_S^2 |S|^2 + \sum_{i=1}^3 (m_{S_i}^2 |S_i|^2 + m_{\tilde{N}_i^c}^2 |\tilde{N}_i^c|^2 + m_{\tilde{\nu}_i}^2 |\tilde{\nu}_i|^2 + m_{\tilde{B}_i}^2 |\tilde{B}_i|^2) - (A_h hSH_1H_2 + A_\lambda \lambda S_1S_2S_3 + A_{d_{ij}} d_{ij}S\tilde{B}_i\tilde{B}_j + A_{e_{ij}} e_{ij}S_3N_i^c\tilde{B}_j + A_{y_{ij}} y_{ij}H_2N_i^cL_j + \text{H.c.}) + (m_{SS_1}^2 SS_1 + m_{SS_2}^2 SS_2 + \text{H.c.}). \quad (39)$$

For simplicity, we do not consider the soft mass terms like $S_1^\dagger S_2$ or $\tilde{N}_i^{c\dagger} \tilde{N}_j^c$ or $\tilde{B}_i^\dagger \tilde{B}_j$, etc.

The runaway directions for the unbounded from below scalar potential are discussed in the appendix, where suitable conditions to avoid them are given. Because we choose relatively large and positive soft mass-squares for $\tilde{\nu}_i$, \tilde{N}_i^c and \tilde{B}_i that are of order 200 GeV or A_h , the scalar fields $\tilde{\nu}_i$, \tilde{N}_i^c and \tilde{B}_i do not acquire nonzero VEVs. Thus, the VEVs for the H_1^0 , H_2^0 , S and S_i are the same as those in [19], and the $Z - Z'$ mass hierarchy and the particle spec-

trum for charginos, neutralinos and Higgs particles are unchanged.

In the basis $\{\nu_1, \nu_2, \nu_3, N_1^c, N_2^c, N_3^c, B_1, B_2, B_3\}$, the neutrino mass matrix is

$$M = \begin{pmatrix} 0 & M_D & 0 \\ M_D^T & 0 & M_V \\ 0 & M_V^T & M_B \end{pmatrix}, \quad (40)$$

where

$$(M_D) = y'_{ij}\nu_2, \quad (M_V) = e_{ij}s_3, \quad (M_B) = d_{ij}s/2, \quad (41)$$

with $\langle H_2^0 \rangle = \nu_2$, $\langle S_3 \rangle = s_3$, $\langle S \rangle = s$, and the upper index T denotes the transpose.

Define the matrix U as

$$U = \begin{pmatrix} 1 & \frac{\sqrt{2}}{2}M_D^*(M_V^{-1})^\dagger & \frac{\sqrt{2}}{2}M_D^*(M_V^{-1})^\dagger \\ (M_V^{-1})^T M_B M_V^{-1} M_D^T & \frac{\sqrt{2}}{2} & -\frac{\sqrt{2}}{2} \\ -M_V^{-1} M_D^T & \frac{\sqrt{2}}{2} & \frac{\sqrt{2}}{2} \end{pmatrix}. \quad (42)$$

$U^T M U$ is approximately [up to $\mathcal{O}(M_V^{-3})$] block diagonal with an upper 3×3 block which gives 3 very light active neutrinos, and a 6×6 block which gives 6 heavy SM singlets. The 3×3 matrix for the active neutrinos is

$$M_\nu = M_D(M_V^{-1})^T M_B M_V^{-1} M_D^T. \quad (43)$$

Using our previous numerical results for the vacuum in Ref. [19], we have that $\nu_2 \sim 125$ GeV, $s \sim 187$ GeV and $s_3 \sim 1260$ GeV. Therefore, realistic active neutrino masses can be obtained, e.g., for y_{ij} and d_{ij} of order 10^{-3} , and e_{ij} of order 1.

We consider the real case for simplicity. The discussions for the complex case are similar. M_D and M_V are general 3×3 mass matrices which have 9 independent parameters, and M_B is a symmetric matrix without diagonal entries, which has 3 independent parameters. However, only $M_D(M_V^{-1})^T$ and M_B enter the expression for M_ν , so, there are $9 + 3 = 12$ independent parameters. Using Mathematica, one can show that M_ν is equivalent to a general real and symmetric mass matrix for the active neutrinos, which has six independent parameters.

It is not hard to find examples which lead to realistic neutrino mass matrices. Here, we consider simple patterns corresponding to a normal hierarchy, an inverted hierarchy with the same signs for the eigenvalues m_{ν_1} and m_{ν_2} , an inverted hierarchy with opposite signs, and the degenerate case.

Define the matrices

$$\alpha = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 1 & 1 \\ 0 & 1 & 1 \end{pmatrix}, \quad \beta = \begin{pmatrix} 2 & 0 & 0 \\ 0 & 1 & -1 \\ 0 & -1 & 1 \end{pmatrix}, \quad (44)$$

$$\gamma = \begin{pmatrix} 0 & 1 & -1 \\ 1 & 0 & 0 \\ -1 & 0 & 0 \end{pmatrix},$$

which correspond to the zeroth order approximations for the patterns of the normal hierarchy, inverted hierarchy with same sign eigenvalues, and inverted hierarchy with opposite sign eigenvalues, respectively. α and β lead to maximal atmospheric neutrino mixing, with the solar neutrino mixing depending on the subleading terms (not displayed) and on the charged lepton mixings. γ leads to bimaximal mixings, which can be consistent with the observed (nonmaximal) solar neutrino mixing if there is small (Cabibbo-like) mixing in the charged lepton sector [41]. Define the mass matrix M'_ν as

$$M'_\nu = X\alpha + Y\beta + Z\gamma. \quad (45)$$

For simplicity, we consider the scenarios in which M_D , M_V and M_B are order unity, i.e., the magnitudes of the entries are $\mathcal{O}(1)$ or $\mathcal{O}(0)$, and show that one can produce the above simple patterns and the patterns with degenerate masses. One can use the freedom in the right-hand side of Eq. (43) to choose

$$M_V = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (46)$$

$$M_B = \begin{pmatrix} 0 & 1 & 1 \\ 1 & 0 & -1 \\ 1 & -1 & 0 \end{pmatrix}, \quad (47)$$

$$M_D = \begin{pmatrix} \sqrt{(d-e)f + d(e+f)} & b & \sqrt{(d-e)f + d(e+f)} \\ d & e & f \\ -d & e & -f \end{pmatrix}. \quad (48)$$

Requiring that $M_\nu = M'_\nu$, we obtain

$$X = (d-f)e, \quad Y = (d-f)e + 2df, \quad (49)$$

$$Z = b(d-f) + (d+f)\sqrt{(d-e)f + d(e+f)}. \quad (50)$$

In the following, we give the solutions for five simple patterns:

- (1) The normal hierarchy: $X \neq 0$ and $Y = Z = 0$. A simple solution to Eqs. (49) and (50) is that $b = 0$ and $d = e = -f = \sqrt{X/2}$, so the Dirac mass matrix M_D is

$$M_D = \sqrt{\frac{X}{2}} \begin{pmatrix} 0 & 0 & 0 \\ 1 & 1 & -1 \\ -1 & 1 & 1 \end{pmatrix}. \quad (51)$$

- (2) The inverted hierarchy with same signs for the eigenvalues m_{ν_1} and m_{ν_2} : $Y \neq 0$ and $X = Z = 0$. A simple real solution is $b = -3\sqrt{Y}$, $d = \sqrt{Y}$, $e = 0$ and $f = \sqrt{Y}/2$, implying

$$M_D = \frac{\sqrt{Y}}{2} \begin{pmatrix} 2 & -6 & 2 \\ 2 & 0 & 1 \\ -2 & 0 & -1 \end{pmatrix}. \quad (52)$$

For M_D complex, there is a simple solution in which $b = 0$, $d = -f = i\sqrt{Y/2}$ and $e = 0$, so,

$$M_D = \sqrt{\frac{Y}{2}} \begin{pmatrix} \sqrt{2} & 0 & \sqrt{2} \\ i & 0 & -i \\ -i & 0 & i \end{pmatrix}. \quad (53)$$

- (3) The inverted hierarchy with opposite signs: $Z \neq 0$ and $X = Y = 0$. A simple solution is $b = d = \sqrt{Z}$, $e = 0$ and $f = 0$, thus,

$$M_D = \sqrt{Z} \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ -1 & 0 & 0 \end{pmatrix}. \quad (54)$$

- (4) The degenerate scenario: $X = Y \neq 0$ and $Z = 0$, can be obtained for $b = -d = -e = -\sqrt{X}$ and $f = 0$, with

$$M_D = \sqrt{X} \begin{pmatrix} 1 & -1 & 1 \\ 1 & 1 & 0 \\ -1 & 1 & 0 \end{pmatrix}. \quad (55)$$

- (5) The degenerate scenario: $X = Z \neq 0$ and $Y = 0$, corresponds to $b = d = e = -f = \sqrt{X/2}$, yielding

$$M_D = \sqrt{\frac{X}{2}} \begin{pmatrix} 0 & 1 & 0 \\ 1 & 1 & -1 \\ -1 & 1 & 1 \end{pmatrix}. \quad (56)$$

V. DISCUSSIONS AND CONCLUSIONS

In this paper, we considered neutrino masses in supersymmetric models with an additional TeV-scale $U(1)'$ gauge symmetry, in which the ordinary seesaw mechanism may not work unless the right-handed neutrinos have no $U(1)'$ charge. We proposed three mechanisms for neutrino masses in such models. First, in models with the gauge group $G_{SM} \times U(1)' \times U(1)''$, with the $U(1)''$ breaking at the intermediate scale, the neutrinos may obtain small Dirac masses through high-dimensional operators associated with the intermediate scale. We illustrated this mechanism in a model with the E_6 particle content and charge assignments, and showed that the right-handed neutrinos could naturally decouple from the TeV-scale $U(1)'$, thus avoiding cosmological and astrophysical constraints. We also discussed this mechanism for models with a secluded $U(1)'$ -breaking sector (in which the $Z - Z'$ mass hierarchy can be generated naturally) and an intermediate-scale $U(1)''$. In this case the right-handed neutrinos are charged under $U(1)'$ unless one goes outside of the E_6 framework for the charge assignments of the SM singlets. We also considered the possibility that the large Majorana masses for right-handed neutrinos can be generated through the intermediate-scale $U(1)''$ breaking, leading to an ordinary seesaw.

In addition, we described two models with pairs of heavy triplets, with masses around 10^{14} GeV and 10^8 GeV, respectively. After the electroweak symmetry breaking, the triplets obtain very small VEVs and can give a reasonable left-handed neutrino Majorana mass matrix. One can instead integrate out the heavy triplets and obtain the small left-handed neutrino Majorana masses from the resulting nonrenormalizable operators. The low energy neutral Higgs potential is the same as that in [19] up to negligible corrections.

We also studied models in which very small neutrino Majorana masses can be obtained by the double-seesaw mechanism. The neutrino Yukawa couplings can be of order 10^{-3} , i.e., the neutrino Dirac masses are comparable to the muon mass. We slightly modified the model in Ref. [19] by introducing three right-handed neutrinos and three SM singlets. Runaway directions can be avoided by imposing suitable conditions on the soft terms. The vacuum is the same as in [19], so the $Z - Z'$ mass hierarchy and the particle spectrum are not modified. The active neutrino mass matrix is $M_D(M_V^{-1})^T M_B M_V^{-1} M_D^T$, where the typical mass scale for M_V is TeV [the $U(1)'$ -breaking scale]. The active neutrinos can have realistic masses and mixings if the typical mass scales for M_B and M_D are about 0.1 GeV. Specific examples for the form of the neutrino Dirac mass matrix that lead to normal and inverted hierarchies and to the degenerate scenario are given.

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APPENDIX

We first consider the scalar potential in Sec. IV without neutrino Yukawa couplings, i.e., in the limit $y_{ij} = 0$. The condition for 6 heavy SM singlets at the TeV-scale is $\det|e_{ij}| \neq 0$. The discussions of the unbounded from below runaway directions for the scalar potential are standard, so, we will not give the details. The constraint conditions to avoid them are

$$m_S^2 + m_{S_1}^2 - 2|m_{SS_1}^2| > 0, \quad (\text{A1})$$

$$m_S^2 + m_{S_2}^2 - 2|m_{SS_2}^2| > 0, \quad (\text{A2})$$

$$\frac{2}{3}m_S^2 + m_{\tilde{N}_i^c}^2 > 0, \quad (\text{A3})$$

$$\frac{3}{2}m_S^2 - \frac{3}{2}|m_{SS_1}^2| + m_{\tilde{N}_i^c}^2 > 0, \quad (\text{A4})$$

$$\frac{3}{2}m_S^2 - \frac{3}{2}|m_{SS_2}^2| + m_{\tilde{N}_i^c}^2 > 0, \quad (\text{A5})$$

where $i = 1, 2, 3$. The constraint condition in Eq. (A1) [or Eq. (A2)] avoids runaway directions in which $\langle S \rangle$ and $\langle S_1 \rangle$ (or $\langle S_2 \rangle$) go to infinity while the other fields have finite VEVs. Equation (A3) avoids the runaway directions in

which $\langle S \rangle$ and $\langle \tilde{N}_i^c \rangle$ (or two or three $\langle \tilde{N}_i^c \rangle$) go to infinity. The condition in Eq. (A4) [or Eq. (A5)] avoids the runaway directions for which $\langle S \rangle$, $\langle S_1 \rangle$ (or $\langle S_2 \rangle$), and $\langle \tilde{N}_i^c \rangle$ (or two or three $\langle \tilde{N}_i^c \rangle$) go to infinity. We assume that $m_{\tilde{N}_i^c}$ are positive, so Eq. (A3) is satisfied automatically if Eq. (A4) or Eq. (A5) is satisfied.

Now let us include the neutrino Yukawa couplings. For simplicity, we assume that $A_{d_{ij}} = A_{e_{ij}} = A_{y_{ij}} = 0$. The only new possible runaway directions have $|\nu_i| \rightarrow \infty$, i.e., $\langle \tilde{\nu}_i \rangle$, $\langle S_3 \rangle$, $\langle \tilde{B}_i \rangle$, $\langle H_1^0 \rangle$ and $\langle H_2^0 \rangle$ can go to infinity, while the other fields have finite VEVs. The VEVs for $\langle \tilde{\nu}_i \rangle$, $\langle S_3 \rangle$, $\langle \tilde{B}_i \rangle$, $\langle H_1^0 \rangle$ and $\langle H_2^0 \rangle$ must satisfy

$$e_{ij}\langle S_3 \rangle \langle \tilde{B}_j \rangle = -h_{ij}\langle H_2^0 \rangle \langle \tilde{\nu}_j \rangle, \quad (\text{A6})$$

$$h\langle H_1^0 \rangle \langle H_2^0 \rangle = -d_{ij}\langle \tilde{B}_i \rangle \langle \tilde{B}_j \rangle, \quad (\text{A7})$$

$$|\langle H_2^0 \rangle|^2 = |\langle H_1^0 \rangle|^2 + \sum_{i=1}^3 |\langle \tilde{\nu}_i \rangle|^2, \quad (\text{A8})$$

$$\begin{aligned} \frac{1}{2} \sum_{i=1}^3 |\langle \tilde{B}_i \rangle|^2 &= 2|\langle S_3 \rangle|^2 + Q_{H_1} |\langle H_1^0 \rangle|^2 + Q_{H_2} |\langle H_2^0 \rangle|^2 \\ &+ \sum_{i=1}^3 Q_L |\langle \tilde{\nu}_i \rangle|^2. \end{aligned} \quad (\text{A9})$$

The potential is very complicated, so we will impose strong conditions to avoid the runaway direction. For the F -terms, we only keep the term $h^2|S|^2|H_2|^2$ because F -terms give positive contributions to the potentials. Using Eqs. (A6)–(A9), we obtain

$$\begin{aligned} V_{\text{total}} &> \left(m_{H_1}^2 + m_{H_2}^2 + \frac{m_{S_3}^2}{2} - A_h^2 \right) |H_1^0|^2 \\ &+ \left(m_{\tilde{B}_i}^2 + \frac{m_{S_3}^2}{4} \right) |\tilde{B}_i|^2 \\ &+ \left(m_{\tilde{\nu}_i}^2 + m_{H_2}^2 + \frac{Q_L + Q_{H_2}}{2} m_{S_3}^2 \right) |\tilde{\nu}_i|^2 \\ &+ \text{constant}, \end{aligned} \quad (\text{A10})$$

where $V_{\text{total}} = V_F + V_D + V_{\text{soft}}$. Because $|\langle H_1^0 \rangle| < |\langle H_2^0 \rangle|$ and we consider the large A_h scenario in which $m_{H_1}^2 + m_{H_2}^2 + \frac{m_{S_3}^2}{2} - A_h^2 < 0$, we obtain

$$\begin{aligned}
V_{\text{total}} &> \left(m_{H_1}^2 + m_{H_2}^2 + \frac{m_{S_3}^2}{2} - A_h^2 \right) \frac{1}{h} |d_{ij}| |\tilde{B}_i| |\tilde{B}_j| \\
&+ \left(m_{\tilde{B}_i}^2 + \frac{m_{S_3}^2}{4} \right) |\tilde{B}_i|^2 \\
&+ \left(m_{\tilde{\nu}_i}^2 + m_{H_2}^2 - \frac{Q_L + Q_{H_2}}{2} m_{S_3}^2 \right) |\tilde{\nu}_i|^2 \\
&+ \text{constant.}
\end{aligned} \tag{A11}$$

$$m_{\tilde{\nu}_i}^2 + m_{H_2}^2 - \frac{Q_L + Q_{H_2}}{2} m_{S_3}^2 > 0, \tag{A12}$$

$$m_{\tilde{B}_i}^2 + \frac{m_{S_3}^2}{4} > \frac{A_h^2 - m_{H_1}^2 - m_{H_2}^2 - \frac{m_{S_3}^2}{2}}{h} \max\{|d_{ij}|, |d_{ik}|\}, \tag{A13}$$

where $i \neq j \neq k$ and $i, j, k = 1, 2, 3$.

Because $m_{\tilde{\nu}_i}^2$, $m_{N_i^c}^2$, and $m_{\tilde{B}_i}^2$ are relatively large positive soft mass-squares of the order of A_h^2 , the above conditions are satisfied.

To avoid the runaway directions, we require

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