A model realizing inverse seesaw and resonant leptogenesis

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We construct a model realizing the inverse seesaw mechanism. The model has two types of gauge singlet fermions in addition to right-handed neutrinos. A required Majorana mass scale (keV scale) for generating the light active neutrino mass in the conventional inverse seesaw can be naturally explained by a "seesaw" mechanism between the two singlet fermions in our model. We find that our model can decrease the magnitude of hierarchy among the mass parameters by $\mathcal{O}(10^4)$ from that in the conventional inverse seesaw model. We also show that a successful resonant leptogenesis occurs for generating the baryon asymmetry of the universe in our model. The desired mass degeneracy for the resonant leptogenesis can also be achieved by the "seesaw" between the two singlet fermions.

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Subject Index   B40, B54

1.   Introduction

Small neutrino masses are observed in neutrino oscillation experiments. One simple mechanism to generate the small neutrino masses is the conventional Type-I seesaw mechanism [1,2,4–7,9], in which right-handed neutrinos N_R^i (i denotes the generation) are introduced to the standard model (SM). The masses of light active neutrinos are described by $M_\nu \simeq -M_D M_R^{-1} M_D^T$ when $|(M_D)_{\alpha i}| \ll |(M_R)_{ij}|$, where α denotes the flavor, M_D is the Dirac neutrino mass matrix given by the Yukawa coupling matrix of neutrinos Y_ν and the vacuum expectation value (VEV) of the Higgs v ($M_D \equiv Y_\nu v$), and M_R is the Majorana mass matrix for N_R. When one takes the magnitude of the neutrino Yukawa couplings as $\mathcal{O}(1)$, like the top Yukawa coupling, the typical size of the right-handed neutrino Majorana mass becomes $M_R \sim \mathcal{O}(10^{14})$ GeV to generate the light active neutrino masses as $m_\nu \sim \mathcal{O}(0.1)$ eV. On the other hand, right-handed neutrinos with electroweak (EW)-scale masses require relatively small M_D of $\mathcal{O}(10^{-4})$ GeV to realize the active neutrino mass scale of $\mathcal{O}(0.1)$ eV in the Type-I seesaw mechanism.

There are several extensions of the Type-I seesaw model. One extension is the inverse (double) seesaw mechanism [8–10] with additional singlet fermions S_α. In a basis of $\left(\nu_L^c, N_R, S\right)^T$ with three flavors (generations), a neutrino mass matrix is given by

$$\mathcal{M} = \begin{pmatrix} 0 & M_D & 0 \\ M_D^T & 0 & M_S \\ 0 & M_S^T & \mu \end{pmatrix}. \tag{1}$$

When one assigns the lepton number one unit to ν_L, N_R, and S (S has an opposite lepton number with respect to that of N_R), the Majorana mass terms of S do not conserve the lepton number. Note that the absence of $(\mathcal{M})_{13}$ and $(\mathcal{M})_{31}$ in Eq. (1) are ensured by a field redefinition. Assuming $\mu \ll M_D < M_S$, one can describe the neutrino mass as $M_\nu \simeq \mu M_D^2 / M_S^2$ by diagonalizing the above mass matrix. For $M_D = 10\,\text{GeV}$ and $M_S = 1\,\text{TeV}$, $\mu \simeq 1\,\text{keV}$ is required for generating the small active neutrino mass scale. In this model, one obtains heavy Majorana neutrinos with masses $M_S \pm \mu/2$. Thus, when $\mu \ll M_S$, such neutrinos are degenerate in mass and can realize the resonant leptogenesis mechanism to generate the baryon asymmetry of the universe (BAU) [11–14]. The explanations of neutrino experimental data and dark matter in the generic class of the inverse seesaw model have been discussed in Refs. [15] and [16], respectively.

In this work, we will discuss the inverse seesaw model realized by a "seesaw" mechanism in the TeV-scale physics. Our model has two kinds of new gauge singlet fermions S_1 and S_2 in addition to N_R, which corresponds to the $n = 2$ multiple seesaw mechanism in Ref. [17]. We will find that our model can naturally induce a very small mass difference between heavy ($\sim$ TeV-scale) neutrino states, which can also be responsible for a successful resonant leptogenesis.

2.   Inverse seesaw from "seesaw"

We discuss a realization of the inverse seesaw from the "seesaw" mechanism. The relevant Lagrangian is given by

$$-\mathcal{L} = Y_\nu \tilde{H} \overline{L} N_R + Y_{S_1} \Phi_1 \overline{N_R^c} S_1 + Y_{S_2} \Phi_2 \overline{S_1^c} S_2 + \frac{M_\mu}{2} \overline{S_2^c} S_2 + h.c., \tag{2}$$

where $\tilde{H} \equiv i\sigma_2 H^*$, H is the SM Higgs doublet, L is the left-handed lepton doublet, N_R are the right-handed neutrinos, Φ_1 and Φ_2 are gauge singlet scalars under the SM gauge groups, S_1 and S_2 are gauge singlet fermions, and M_μ is the Majorana mass of S_2. Note that S_2 is added to the original inverse seesaw mechanism. Here, details of additional symmetries in our model are not specified, but discussed later. In order to reproduce two (solar and atmospheric) mass scales of the active neutrinos, one must introduce at least two generations for the right-handed neutrino or the gauge singlet fermions. We omit the generation and flavor indices for fermions in Eq. (2). After spontaneous gauge symmetry breaking, one can describe a neutrino mass matrix as

$$\mathcal{M} = \begin{pmatrix} 0 & M_D & 0 & 0 \\ M_D^T & 0 & M_{S_1} & 0 \\ 0 & M_{S_1}^T & 0 & M_{S_2} \\ 0 & 0 & M_{S_2}^T & M_\mu \end{pmatrix}, \tag{3}$$

in the basis of $\left(\nu_L^c, N_R, S_1, S_2\right)^T$, where $M_D \equiv Y_\nu \langle H \rangle$, $M_{S_1} \equiv Y_{S_1} \langle \Phi_1 \rangle$, $M_{S_2} \equiv Y_{S_2} \langle \Phi_2 \rangle$, and these are described by matrices.[1] If one adds three generations for each singlet fermion, the neutrino mass matrix $\mathcal{M}$ is a 12×12 matrix.

When we assume that the values of all the matrix elements of M_{S_2} are much smaller than those of M_μ ($(M_{S_2})_{jk} \ll (M_\mu)_{lm}$), we can diagonalize the lower-right 2×2 sub-matrix (integrate out the S_2

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[1] A similar structure for the mass matrix has been discussed in the three active and two sterile neutrinos model for the liquid scintillator neutrino detector anomaly [18].

field), i.e., utilize a "seesaw" mechanism. Then, block diagonalization gives

$$\mathcal{M} \to \begin{pmatrix} \bar{\mathcal{M}} & 0 \\ 0 & \bar{M}_\mu \end{pmatrix}, \tag{4}$$

with

$$\bar{\mathcal{M}} \simeq \begin{pmatrix} 0 & M_D & 0 \\ M_D^T & 0 & M_{S_1} \\ 0 & M_{S_1}^T & \mu \end{pmatrix}, \quad \bar{M}_\mu \simeq M_\mu, \ \mu \simeq -M_{S_2} M_\mu^{-1} M_{S_2}^T, \tag{5}$$

in a basis. Notice that $\bar{\mathcal{M}}$ takes the same form of the original inverse seesaw as in Eq. (1), and the smallness of μ can be naturally realized by the "seesaw" mechanism, $\mu \simeq -M_{S_2} M_\mu^{-1} M_{S_2}^T$.

The mass matrix for the three active neutrinos M_ν can be obtained after the inverse seesaw as

$$M_\nu \simeq -M_D \left(M_{S_1}^T\right)^{-1} M_{S_2} M_\mu^{-1} M_{S_2}^T \left(M_{S_1}\right)^{-1} M_D^T, \tag{6}$$

in the flavor basis of active neutrinos. The magnitude of the matrix elements of the active neutrino is realized as

$$M_\nu \sim \left(\frac{M_D}{10\,\text{GeV}}\right)^2 \left(\frac{M_{S_1}}{1\,\text{TeV}}\right)^{-2} \left(\frac{M_{S_2}}{30\,\text{MeV}}\right)^2 \left(\frac{M_\mu}{1\,\text{TeV}}\right)^{-1} 0.1\,\text{eV}. \tag{7}$$

In the conventional (TeV-scale) inverse seesaw mechanism, one should require μ in the mass matrix of Eq. (1) of the $\mathcal{O}(1)$ keV scale. On the other hand, the realization of the inverse seesaw from "seesaw", Eq. (3), needs M_{S_2} of the $\mathcal{O}(10)$ MeV scale instead of keV scale. Therefore, the mass hierarchy among the singlet fermions in our model becomes small compared to the usual inverse seesaw model. Considering that light quarks and leptons have MeV-scale masses, the scale could be usable as a parameter of the model.

3. Leptogenesis

Next, we discuss the generation of the BAU. Our model includes several singlet Majorana fermions, and the masses of some of them can be taken as $\mathcal{O}(1)$ TeV. Thus, resonant leptogenesis [19] might be possible in the model.

We start from the mass matrix, Eq. (5). Since the typical size of the matrix elements of M_{S_1} is much larger than that of μ, the mixing angle for block diagonalization of the lower-right 2×2 sub-matrix of Eq. (5) is almost maximal. Thus, $\bar{\mathcal{M}}$ is rotated as

$$\bar{\mathcal{M}} \to \bar{\mathcal{M}}' \simeq \begin{pmatrix} 0 & M_D(1+\epsilon)/\sqrt{2} & M_D(1+\epsilon)/\sqrt{2} \\ (M_D(1+\epsilon))^T/\sqrt{2} & M_{S_1} - \mu/2 & 0 \\ (M_D(1+\epsilon))^T/\sqrt{2} & 0 & M_{S_1} + \mu/2 \end{pmatrix}, \tag{8}$$

up to the order of $\mathcal{O}(\mu)$ in a basis of $\left(\nu_L^c, X_-, X_+\right)^T$, where $\epsilon \simeq \mu/\left(2\sqrt{2}M_{S_1}\right)$, and the eigenstates $X_\pm$ are described as $X_\pm \simeq \left(c_R N_R \mp c_{1R} S_1 \mp c_1 S_2\right)/\sqrt{2}$ with $c_R \simeq c_{1R} \simeq 1$, where c_1 is estimated as a typical ratio of matrix elements of M_{S_2} and M_μ, $c_1 \simeq \mathcal{O}\left(M_{S_2}/M_\mu\right)$. The relevant Lagrangian for the resonant leptogenesis is given from Eq. (2) as

$$-\mathcal{L} \supset Y_\nu^N \tilde{H} \overline{L} X_- + Y_\nu^S \tilde{H} \overline{L} X_+ + \frac{M_{S_1} - \mu/2}{2} \overline{X_-^c} X_- + \frac{M_{S_1} + \mu/2}{2} \overline{X_+^c} X_+ + h.c., \tag{9}$$

for the eigenstates of $X_\pm$. Note that the typical size of the matrix elements of Y_ν^N and Y_ν^S is of the same order as that of Y_ν in Eq. (2), $Y_\nu^N \simeq Y_\nu^S \simeq (c_R/\sqrt{2})Y_\nu$ at the leading order. Hereafter, we assume that the 3×3 matrices M_{S_1} and μ are diagonal matrices, for simplicity. We also assume a hierarchical structure for M_{S_1} as $m_{S_1} \equiv (M_{S_1})_{11} \ll (M_{S_1})_{22}, (M_{S_1})_{33}$ so that the BAU can be induced by the decays of the first generation of $X_\pm$ ($\equiv \chi_\pm$), whose masses are obtained as $m_{\chi_\pm} = m_{S_1} \pm \mu/2$.[2]

The lepton asymmetry from the decays of χ_- and χ_+ is calculated as [11,24]

$$
\epsilon_\pm = \frac{\sum_\alpha \left[\Gamma(\chi_\pm \to L_\alpha + H^*) - \Gamma(\chi_\pm \to L_\alpha^c + H) \right]}{\sum_\alpha \left[\Gamma(\chi_\pm \to L_\alpha + H^*) + \Gamma(\chi_\pm \to L_\alpha^c + H) \right]}
$$
$$
\simeq \frac{\mathrm{Im} \left(Y_\nu^{N\dagger} Y_\nu^S Y_\nu^{N\dagger} Y_\nu^S \right)_{11}}{8\pi A_\pm} \frac{r}{r^2 + \Gamma_\mp^2/m_{\chi_\mp}^2}, \tag{10}
$$

where

$$
r = \frac{m_{\chi_+}^2 - m_{\chi_-}^2}{m_{\chi_+} m_{\chi_-}} \simeq \frac{2\mu}{m_{S_1}}, \quad A_+ = \left(Y_\nu^{S\dagger} Y_\nu^S \right)_{11}, \quad A_- = \left(Y_\nu^{N\dagger} Y_\nu^N \right)_{11}, \tag{11}
$$

and $\Gamma_\pm = A_\pm m_{\chi_\pm}/(8\pi)$ is the decay width of $\chi_\pm$. The baryon asymmetry is given by the lepton asymmetry as

$$
\eta_B = -\frac{28}{79} \frac{0.3\epsilon_\pm}{g_* K_\pm (\ln K_\pm)^{0.6}}, \tag{12}
$$

where $g_* = 106.75$ is the relative degree of freedom and $K_\pm = \Gamma_\pm/(2H(T))|_{T=m_{\chi_\pm}}$ with the Hubble constant $H(T) = 1.66\sqrt{g_*}T^2/m_{\mathrm{Pl}}$. Note that the baryon asymmetry is enhanced for $(m_{\chi_+} - m_{\chi_-}) \sim \Gamma_\pm/2$.

In order to obtain the baryon asymmetry by the decays of $\chi_\pm$, the $\chi_\pm$ should be decoupled at $T \sim m_{\chi_\pm}$, which is realized for the Yukawa couplings $(Y_\nu^N)_{\alpha 1}$ and $(Y_\nu^S)_{\alpha 1}$, being $< \mathcal{O}(10^{-6})$. Under these conditions, the appropriate order of r in Eq. (11) for the resonant leptogenesis is $r \sim 10^{-9}$, which can also be naturally realized in our model. Regarding the masses of additional scalars $\Phi_{1,2}$, these must be larger than the masses of $\chi_\pm$ of $\mathcal{O}(1)$ TeV. If these masses are smaller than the TeV scale, $\chi_\pm$ decay into scalars. As a result, the lepton asymmetry cannot be produced. On the other hand, the VEV of Φ_2 should be larger than $\mathcal{O}(10)$ MeV to realize the inverse seesaw when $Y_{S_2} \leq \mathcal{O}(1)$. Such a hierarchy between the mass and VEV can be realized in the neutrino-philic Higgs model [25] (see also Refs. [26,27]). The relevant scalar potential for the realization is, e.g.,

$$
V \supset -m_{\Phi_1}^2 |\Phi_1|^2 + m_{\Phi_2}^2 |\Phi_2|^2 - m^2 \left(\Phi_1^* \Phi_2 + \Phi_2^* \Phi_1 \right) + \frac{\lambda_{\Phi_1}}{2} |\Phi_1|^4 + \frac{\lambda_{\Phi_2}}{2} |\Phi_2|^4
$$
$$
+ \lambda_3 |\Phi_1|^2 |\Phi_2|^2 + \lambda_4 |\Phi_1 \Phi_2|^2 + \frac{\lambda_5}{2} \left[\left(\Phi_1^* \Phi_2 \right)^2 + \left(\Phi_2^* \Phi_1 \right)^2 \right], \tag{13}
$$

where $m_{\Phi_{1,2}}$, m, $\lambda_{\Phi_{1,2}}$, and $\lambda_{3,4,5}$ are all assumed to be real and positive, for simplicity. The stationary conditions $\partial V/\partial\langle\Phi_1\rangle = 0$ and $\partial V/\partial\langle\Phi_2\rangle = 0$ lead to $|\langle\Phi_1\rangle| \simeq m_{\Phi_1}/\sqrt{\lambda_{\Phi_1}}$ and $|\langle\Phi_2\rangle| \simeq m^2|\langle\Phi_1\rangle|/m_{\Phi_2}^2$, respectively, where we assume $\lambda_{\Phi_1}, \lambda_{3,4,5} \ll \mathcal{O}(1)$ and $\lambda_{\Phi_2}|\langle\Phi_2\rangle|^2 \ll m_{\Phi_2}^2$. In addition, when one introduces the symmetry, which forbids the term $m^2 \left(\Phi_1^* \Phi_2 + \Phi_1 \Phi_2^* \right)$ in Eq. (13), the hierarchy $m \ll m_{\Phi_{1,2}}$ seems to be natural. As a result, $|\langle\Phi_2\rangle| \ll |\langle\Phi_1\rangle|$ can be realized, where an assumption $\lambda_{\Phi_2}|\langle\Phi_2\rangle|^2 \ll m_{\Phi_2}^2$ is consistent with this realization. Thus, one can have a hierarchy

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[2] Such a hierarchical mass structure among singlet Majorana fermions can be realized by several models [20–23].
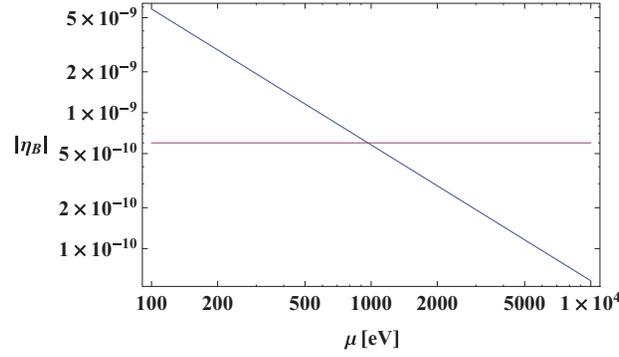
Fig. 1. Baryon asymmetry as a function of μ.

between the VEVs, $|\langle\Phi_1\rangle| \simeq m_{\Phi_1}/\sqrt{\lambda_{\Phi_1}} \gtrsim \mathcal{O}(1)$ TeV and $\langle\Phi_2\rangle = \mathcal{O}(10)$ MeV, when one takes $m = \mathcal{O}(10)$ GeV and the masses of $\Phi_{1,2}$ as $\mathcal{O}(1)$ TeV. Here we assume that the masses of $\Phi_{1,2}$ are larger than the masses of $\chi_\pm$. The above calculation is valid in this case. In our model, the other singlet fermions $(\simeq S_2)$ can also decay into L and H through the mixing between S_1 and S_2. However, the process cannot generate a sufficient magnitude of lepton asymmetry because S_2 is not degenerate with the N_R and S_1 states. Note that one does not need fine tuning between the masses of Majorana fermions to realize the BAU, as seen below.

Figure 1 shows the baryon asymmetry as a function of μ. We assume hierarchical structures for the Yukawa couplings Y_ν^N and Y_ν^S so as to realize the nature of $\chi_\pm$ as being out of thermal equilibrium at $T \sim m_{\chi\pm}$: $\left|(Y_\nu^N)_{\alpha1}\right|, \left|(Y_\nu^S)_{\beta1}\right| < \mathcal{O}(10^{-6})$. The observed baryon asymmetry $\eta_B = 6 \times 10^{-10}$ is also shown by the horizontal line in Fig. 1. In the calculation, we take $(Y_\nu^N)_{\alpha1} = (1.0 + 0.1\,i) \times 10^{-6}$, $(Y_\nu^S)_{\alpha1} = (1.0 + 0.3\,i) \times 10^{-6}$, and $M_{S_1} = 1$ TeV as reference values. It is seen that the observed baryon asymmetry can be realized at $\mu \simeq 1$ keV. The case of $\mu \simeq \mathcal{O}(1)$ keV is consistent with the realization of the small active neutrino mass in our model, i.e. $\mu \simeq M_{S_2} M_\mu^{-1} M_{S_2}^T \simeq \mathcal{O}(1)$ keV.

As discussed above, the favored scale of μ for the active neutrino mass can be realized by the "seesaw" between $S_{1,2}$ fermions. In addition, the first generations of S_1 and N_R (those mass eigenstates are $\chi_\pm$) play a role in generating the BAU via resonant leptogenesis. In this case, the required size of the mass degeneracy between $\chi_\pm$ for the resonant leptogenesis, $\mu \simeq \mathcal{O}(1)$ keV, can also be realized by the "seesaw" between $S_{1,2}$. Both realizations are non-trivial results in our model.

4. Signatures for the LHC experiment

We discuss signatures of this model in the LHC experiment. This model can induce lepton-number-violating processes. One interesting process is like-sign dilepton production, $q\bar{q}' \to \ell^\pm\ell^\pm W^\mp$, where the lepton-number conservation is violated by two units, $\Delta L = 2$, due to the Majorana nature of the neutrinos. References [28–35] explore this process at the LHC experiment in the SM with right-handed Majorana neutrinos (see also Ref. [36] for a review of the collider phenomenology with right-handed and sterile Majorana neutrinos).[3] According to Ref. [31], it is found that there is 2σ (5σ) sensitivity for the $\mu^\pm\mu^\pm$ modes in the mass range of a Majorana neutrino of 10 GeV $\leq m_\chi \leq 350$ (250) GeV at the 14 TeV LHC experiment with 100 fb^{-1}. Regarding the

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[3] An analysis of this process is given in Ref. [37] for the inverse seesaw model in the context of the next-to-minimal supersymmetric SM.

inverse seesaw case, the singlet neutrinos and fermions are pseudo-Dirac neutrinos due to a small Majorana mass μ, and the neutrinos contain tiny Majorana states. The ratio of the Majorana state is typically determined by $\mu/M_{S_1} \simeq 1\,\text{keV}/1\,\text{TeV} \simeq \mathcal{O}(10^{-9})$. Thus, since the like-sign dilepton production process in the inverse seesaw case is suppressed by $(\mu/M_{S_1})^2 \simeq \mathcal{O}(10^{-18})$ compared with the results of Refs. [28–31], the signatures of the process in the inverse seesaw case cannot reach the sensitivity of the LHC experiment.

Similarly, for other singlet fermions ($\simeq S_2$) with the lepton-number-violating Majorana mass of $\mathcal{O}(1)$ TeV in our model, the result of the analysis in Refs. [28–31] cannot simply be adopted. Since the like-sign dilepton production process is induced through the mixing between S_1 and S_2 in addition to the mixing of the pseudo-Dirac states of N_R and S_1 mentioned above, the amplitude is typically suppressed by $(M_{S_2}/M_\mu)^2(\mu/M_{S_1})^2 \simeq (10\,\text{MeV}/1\,\text{TeV})^2(1\,\text{keV}/1\,\text{TeV})^2 \simeq \mathcal{O}(10^{-28})$. Therefore, the collider signatures of these singlet fermions in our model also cannot reach the sensitivity of the LHC experiment.

The above discussion can be generalized to multiple seesaw models [17]. For the $n = 2k + 1$ ($k = 0, 1, 2, \ldots$) multiple seesaw models ($k = 0$ is the conventional inverse seesaw model), the active neutrino mass matrix in the $n = 2k + 1$ ($k \geq 1$) multiple seesaw models is given by

$$M_\nu = M_D \left[\prod_{i=1}^{k} \left(M_{S_{2i-1}}^T \right)^{-1} M_{S_{2i}} \right] \left(M_{S_{2k+1}}^T \right)^{-1} M_\mu \left(M_{S_{2k+1}} \right)^{-1} \left[\prod_{i=1}^{k} \left(M_{S_{2i-1}}^T \right)^{-1} M_{S_{2i}} \right]^T M_D^T,$$
(14)

where n denotes the number of gauge singlet fermions S without the number of generation (flavor) and M_μ is the lower-right element of the $(n + 2) \times (n + 2)$ generalized neutrino mass matrix. The like-sign dilepton production process is suppressed by $(M_\mu/M_{S_n})^2$ in all models of $n = 2k + 1$ multiple seesaw with $M_\mu \simeq \mathcal{O}(1)\,\text{keV} \ll M_{S_1} \simeq \cdots \simeq M_{S_{n-1}}$. On the other hand, for the $n = 2k$ ($k = 1, 2, \ldots$) multiple seesaw models (the $k = 1$ case is our model), the active neutrino mass matrix can be given by

$$M_\nu = -M_D \left[\prod_{i=1}^{k} \left(M_{S_{2i-1}}^T \right)^{-1} M_{S_{2i}} \right] M_\mu^{-1} \left[\prod_{i=1}^{k} \left(M_{S_{2i-1}}^T \right)^{-1} M_{S_{2i}} \right]^T M_D^T.$$
(15)

The amplitude of the like-sign dilepton production process is suppressed by $\left(M_{S_n}/M_\mu \right)^2 \times (1\,\text{keV}/M_{S_1})^2$ in all models of $n = 2k$ multiple seesaw with $M_{S_n} \ll M_\mu \simeq M_{S_1} \simeq \cdots \simeq M_{S_{n-1}}$. Note that, since the $n = 2k$ multiple seesaw model is reduced to the inverse seesaw model, there is an additional suppression $(M_{S_n}/M_\mu)^2$ in the $n = 2k$ cases compared with the $n = 2k + 1$ multiple seesaw models.[4]

5. Summary

We have discussed the inverse seesaw model realized by a "seesaw" mechanism. The conventional inverse seesaw model requires a lepton-number-violating Majorana mass of $\mu \simeq \mathcal{O}(1)$ keV to achieve the light active neutrino mass scale when the Dirac masses are taken as $M_D = 10\,\text{GeV}$ and $M_S = 1\,\text{TeV}$ (see Eq. (1)). The hierarchy among mass scales in the conventional inverse

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[4] In Refs. [38–40], the authors discussed the Higgs signatures via large Yukawa couplings in the inverse seesaw model at the LHC. The second and third generations of $X_\pm$ in our model might be adopted in the discussion, although the first-generation Yukawa couplings are too small to lead to a sufficient signal magnitude.

Table 1. Charge assignment [17] for our model where $\omega_6 \equiv e^{i\pi/3}$.

	L	H	N_R	S_{1R}	S_{2R}	Φ_1	Φ_2
$U(1)_L$	$+1$	0	$+1$	-1	$+1$	0	0
$\mathbb{Z}_6$	ω_6	$+1$	ω_6	ω_6^2	ω_6^3	ω_6^3	ω_6

seesaw model is given by $M_S/\mu \simeq \mathcal{O}(10^9)$. On the other hand, in our model, the Majorana mass is $M_\mu \simeq \mathcal{O}(1)$ TeV for the Dirac masses of $M_D = 10\,\mathrm{GeV}$, $M_{S_1} = 1$ TeV, and $M_{S_2} = 30\,\mathrm{MeV}$ (see Eqs. (3) and (7)). Thus, the magnitude of mass hierarchy in the model can be decreased to $M_\mu/M_{S_2} \simeq \mathcal{O}(10^5)$, which is due to the "seesaw" mechanism between S_1 and S_2 singlet fermions.

We have also considered a leptogenesis scenario with a mass degeneracy for generating the BAU, the so-called resonant leptogenesis. The scenario can be realized by the keV-scale mass degeneracy between the first generations of the right-handed neutrino and one of the singlet fermions. We have shown that such mass degeneracy can also be realized by the "seesaw" in our model, and thus successful resonant leptogenesis is achieved. Regarding the signatures of $q\bar{q}' \to \ell^\pm \ell^\pm W^\mp$ processes in the LHC experiment, our model cannot reach the sensitivity of the LHC due to significant suppression by mixings between the singlet fermions.

Finally, we comment on a realization of our model. One simple way to obtain our model is to introduce a symmetry. In Ref. [17], the global $U(1) \times \mathbb{Z}_{2N}$ symmetry for realizing multiple seesaw models was discussed. Following that, our model (the $n = 2$ multiple seesaw model) can be obtained by imposing the global $U(1) \times \mathbb{Z}_6$ symmetry. Here, the global $U(1)$ symmetry is identified with the lepton number, $U(1)_L$, and a charge assignment under the symmetry is given in Table 1. Note that the Majorana mass term $(M_\mu/2)\overline{S_2^c}S_2$ induces lepton-number violation.

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