Extended double seesaw model for neutrino mass spectrum
and low scale leptogenesis

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Abstract

We consider a variant of seesaw mechanism by introducing extra singlet neutrinos, with which we show how the low scale leptogenesis is realized without imposing the tiny mass splitting between two heavy Majorana neutrinos required in the resonant leptogenesis. Thus, we can avoid the so-called gravitino problem when our scenario is supersymmetrized. We show that an introduction of the new singlet fermion leads to a new contribution which can enhance the lepton asymmetry for certain range of parameter space. We also examine how both the light neutrino mass spectrum and relatively light sterile neutrinos of order a few 100 MeV can be achieved without being in conflict with the constraints on the mixing between the active and sterile neutrinos.

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There exist quite many motivations for postulating the existence of weak interaction singlets such as singlet neutrinos which may arise in most extensions of the standard model (SM); It is by now well known that the smallness of neutrino masses can be understood by introducing heavy singlet Majorana neutrinos via the so-called seesaw mechanism \[1\]. It was also noticed that sterile neutrinos could be a good viable candidate for cosmological dark matter \[2\]. The controversial result from the LSND experiment \[3\] seems to indicate that the complete picture of neutrino sector has need additional neutrinos beyond the three ordinary ones, which must be sterile. The sterile neutrino states can mix with the active neutrinos and such admixtures contribute to various processes which are forbidden in the SM, and affect the interpretations of cosmological and astrophysical observations. Thus, the masses of the sterile neutrinos and their mixing with the active neutrinos are subject to various experimental bounds as well as cosmological and astrophysical constraints.

In addition to the accomplishment of smallness for neutrino masses, another virtue of the seesaw mechanism is that it can be responsible for baryon asymmetry of our Universe via so-called leptogenesis \[4\]. However, the relevant scale of the typical leptogenesis with hierarchical heavy Majorana neutrino mass spectrum is larger than \(\sim 10^9\) GeV \[5\], which makes it impossible to probe the seesaw model at collider experiments in a foreseeable future. And thermal leptogenesis working at such a high mass scale encounters gravitino problem in the framework of supersymmetric SM as well. Thus, a low scale seesaw is desirable to reconcile such problems. We may emphasize that it is also quite nontrivial to naturally realize the hierarchy of the neutrino mass spectrum in the seesaw model.

Resonant leptogenesis with very tiny mass splitting \((M_2 - M_1)/(M_2 + M_1) \sim 10^{-6})\) of heavy Majorana neutrinos with \(M_1 \sim 1\) TeV has been proposed as a successful scenario for a low scale leptogenesis \[6\]. However, such a very tiny mass splitting may appears somewhat unnatural due to the required severe fine-tuning. In order to remedy such problems mentioned above in this letter, we consider a variant of the seesaw mechanism, in which an equal number of gauge singlet neutrinos is introduced on top of the heavy right-handed neutrinos. The model we consider looks similar to the so-called double seesaw mechanism \[7\], but is obviously different from it because mass terms of the heavy right-handed neutrinos, which are not present in the typical double seesaw model, are here introduced \[8\]. Such additional sterile neutrinos can be naturally incorporated into the superstring \(E_6\) \[7\] and/or flipped \(SU(5)\) GUT \[8\]. Unlike to the typical double seesaw model, as will be shown later, our model permit both tiny neutrino masses and relatively light sterile neutrinos of order MeV.
which may play important roles in cosmology and astrophysics. It can also accommodate very tiny mixing between the active and sterile neutrinos demanded from the cosmological and astrophysical observations. Note that the effects of the heavy right-handed neutrinos are cancelled in our seesaw model for the tiny neutrino masses, whereas they may play an important role in generating MeV sterile neutrinos. In particular, we will show that a low scale thermal leptogenesis can be naturally achieved without encountering the problems described above. And we will examine how both the light neutrino mass spectrum and the sterile neutrino mass of order a few 100 MeV can be obtained without being in conflict with the constraints on the mixing between the active and sterile neutrinos. In this scenario, the mass terms of the newly introduced singlet neutrinos are responsible for the hierarchy of the light neutrino mass spectrum. Thus, the structure of the light neutrino mass matrix may be determined by the mechanism of screening of the Dirac flavor structure [9] in our model.

The Lagrangian we propose is given in the charged lepton basis as

$$\mathcal{L} = M_R N_i^T N_i + Y_{D_{ij}} \bar{\nu}_i \phi N_j + Y_{S_{ij}} \bar{N}_i \Psi S_j - \mu_{ij} S_i^T S_j + h.c.,$$

(1)

where $\nu_i, N_i, S_i$ stand for SU(2)$_L$ doublet, right-handed singlet and newly introduced singlet neutrinos, respectively. And $\phi$ and $\Psi$ denote the SU(2)$_L$ doublet and singlet Higgs sectors. The neutrino mass matrix after the scalar fields get vacuum expectation values becomes

$$M_\nu = \begin{pmatrix} 0 & m_{D_{ij}} & 0 \\ m_{D_{ij}} & M_{R_{ii}} & M_{ij} \\ 0 & M_{ij} & -\mu_{ij} \end{pmatrix},$$

(2)

where $m_{D_{ij}} = Y_{D_{ij}} < \phi >, M_{ij} = Y_{S_{ij}} < \Psi >$. Here we assume that $M_R > M \gg \mu, m_D$. After integrating out the right-handed heavy neutrino sector $N_R$ in the above Lagrangian, we obtain the following effective lagrangian,

$$-\mathcal{L}_{eff} = \frac{(m_D^2)_{ij}}{4M_R} \nu_i^T \nu_j + \frac{m_{D_{ik}} M_{kj}}{4M_R} (\bar{\nu}_i S_j + \bar{S}_i \nu_j) + \frac{M_{ij}^2}{4M_R} S_i^T S_j + \mu_{ij} S_i^T S_j.$$  

(3)

After block diagonalization of the effective mass terms in $\mathcal{L}_{eff}$, the light neutrino mass matrix and mixing between the active and sterile neutrinos are given by

$$m_\nu \simeq \frac{1}{2} \frac{m_D}{M} \mu \left( \frac{m_D}{M} \right)^T,$$

(4)

$$\tan 2\theta_s = \frac{2m_D M}{M^2 + 4\mu M_R - m_D^2},$$

(5)
where we omitted the indices of the mass matrices, $m_D, M, M_R$ and $\mu$ for simplicity.

We note that the term $m_D^2/M_R$ corresponding to typical seesaw (type I) mass is cancelled out. On the other hand, the sterile neutrino mass is approximately given as

$$m_s \simeq \mu + \frac{M^2}{4M_R}. \quad (6)$$

Depending on the relative sizes among $M, M_R, \mu$, the mixing angles $\theta_s$ and the sterile neutrino mass $m_s$ are approximately given by

$$\tan 2\theta_s \simeq \sin 2\theta_s \simeq \begin{cases} \frac{2m_D}{M} & \text{(for } M^2 > 4\mu M_R : \text{ Case A}), \\ \frac{m_D}{2\mu M_R} & \text{(for } M^2 \simeq 4\mu M_R : \text{ Case B}), \\ \frac{M^2}{4M_R} & \text{(Case A)}, \\ 2\mu & \text{(Case B)}, \\ \mu & \text{(Case C)} \end{cases} \quad (7)$$

We can see from Eq. (8) that for $M^2 \leq 4\mu M_R$, the size of $\mu$ is mainly responsible for the value of $m_s$. We also notice that the value of the mixing angle $\theta_s$ is suppressed by the scale of $M$ or $M_R$. Thus, very small mixing angle $\theta_s$ can be naturally achieved in our seesaw mechanism. As expected, for Case A and Case B, the mixing angle $\theta_s$ constrained by the cosmological and astrophysical observation as well as laboratory experiments leads to constraints on the size of the ratio $m_\nu/\mu$ through the Eq. (4)

The existence of a relatively light sterile neutrino has nontrivial observable consequences for cosmology and astrophysics, so that the masses of sterile neutrinos and their mixing with the active ones must be subject to the cosmological and astrophysical bounds. There are also some laboratory bounds which are typically much weaker than the astrophysical and cosmological ones. Those bounds turn out to be useful in the case that the cosmological and astrophysical limits become inapplicable. In the light of laboratory experimental as well as cosmological and astrophysical observations, there exist two interesting mass ranges of the sterile neutrinos, one is of order keV, and the other of order MeV. For those mass ranges of sterile neutrinos, their mixing with the active ones are constrained by various cosmological and astrophysical observations as described below.

It has been shown that a sterile neutrino with the mass in keV range appears to be a viable “warm” dark matter candidate. The small mixing angle ($\sin \theta \sim 10^{-6} - 10^{-4}$) between sterile and active neutrinos ensures that sterile neutrinos were never in thermal equilibrium
in the early Universe and this allows their abundance to be smaller than the predictions in thermal equilibrium. Moreover, a sterile neutrino with these parameters is important for the physics of supernova, which can explain the pulsar kick velocities \[13\]. In addition, there are some bounds on the mass of sterile neutrino from the possibility to observe sterile neutrino radiative decays from X-ray observations and Lyman \(\alpha\)-forest observations of order of a few keV.

On the other hand, there exists high mass region \(m_s \gtrsim 100\) MeV of the sterile neutrinos, which is restricted by the CMB bound, meson decays and SN1987A cooling. For this mass range of sterile neutrino, the mixing between the active and sterile neutrinos is constrained to be approximately \(\sin^2 \theta_s \lesssim 10^{-9}\) \[11\]. Such a high mass region may be very interesting, because induced contributions to the active neutrino mass matrix due to mixing between the active and sterile neutrinos can be dominant and they can be responsible for peculiar properties of the lepton mixing such as tri-bimaximal mixing and neutrino mass spectrum \[11\]. In addition, sterile neutrinos with mass 1-100 MeV can be a dark matter candidate for the explanation of the excess flux of 511 keV photons if \(\sin^2 2\theta_s \lesssim 10^{-17}\) \[14\]. The detailed constraints coming from current astrophysical, cosmological and laboratory observations are presented in the figures of Ref. \[11\]. In this letter, we will focus on sterile neutrinos with masses in the range of MeV. Similarly, we can realize keV sterile neutrinos but at the price of naturalness for the magnitudes of some parameters.

First of all, let us consider how low scale leptogenesis can be successfully achieved in our scenario without severe fine-tuning such as very tiny mass splitting between two heavy Majorana neutrinos. In our scenario, the successful leptogenesis can be achieved by the decay of the lightest right-handed Majorana neutrino before the scalar fields get vacuum expectation values. As will be discussed, when our scenario is supersymmetrized, we can escape the so-called gravitino problem because rather light right-handed Majorana neutrino masses are possible in our scenario. In particular, there is a new contribution to the lepton asymmetry which is mediated by the extra singlet neutrinos.

Without loss of generality, we can rotate and rephase the fields to make the mass matrices \(M_{Rij}\) and \(\mu_{ij}\) real and diagonal. In this basis, the elements of \(Y_D\) and \(Y_S\) are in general complex. The lepton number asymmetry required for baryogenesis is given by

\[
\varepsilon_1 = -\sum_i \left[ \frac{\Gamma(N_1 \rightarrow \bar{l}_i H_u) - \Gamma(N_1 \rightarrow l_i H_u)}{\Gamma_{\text{tot}}(N_1)} \right],
\]

where \(N_1\) is the lightest right-handed neutrino and \(\Gamma_{\text{tot}}(N_1)\) is the total decay rate. In
our scenario, the introduction of the new singlet fermion leads to a new contribution to $\varepsilon_1$ which for certain range of parameters can enhance the effect. We show that as a result of this enhancement, even for $M_1$ as low as a few TeV, we can have successful leptogenesis while keeping tiny masses for neutrinos. Therefore, thermal production of $N_1$ does not need too high reheating temperature and the Universe would not encounter the gravitino overproduction \[15\].

**Fig. 1:** Diagrams contributing to lepton asymmetry.

Fig. 1 shows the structure of the diagrams contributing to $\varepsilon_1$. In addition to the diagrams of the standard leptogenesis scenario [16], there is a new contribution of the diagram which corresponds to the self energy correction of the vertex arisen due to the new Yukawa couplings with singlet neutrinos and Higgs sectors. Assuming that the masses of the Higgs sectors and extra singlet neutrinos are much smaller compared to that of the right-handed neutrino, to leading order, we have

$$\Gamma_{\text{tot}}(N_i) = \frac{1}{4\pi} \left( Y_{\nu}^* Y_{\nu}^T + Y_s^* Y_s^T \right)_{ii} M_{R_i}^2$$

so that

$$\varepsilon_1 = \frac{1}{8\pi} \sum_{k \neq 1} \left( [g_V(x_k) + g_S(x_k)] T_{k1} + g_S(x_k) S_{k1} \right),$$

where $g_V(x) = \sqrt{x} \{1 - (1+x)\ln[(1+x)/x]\}$, $g_S(x) = \sqrt{x_k/(1-x_k)}$ with $x_k = M_{R_k}^2 / M_{R_1}^2$ for $k \neq 1$,

$$T_{k1} = \frac{\text{Im}\left( (Y_{\nu} Y_{\nu}^T)_{k1}^2 \right)}{(Y_{\nu} Y_{\nu}^T + Y_s Y_s^T)_{11}}$$

and

$$S_{k1} = \frac{\text{Im}\left( (Y_{\nu} Y_{\nu}^T)_{k1} (Y_s Y_s^T)_{1k} \right)}{(Y_{\nu} Y_{\nu}^T + Y_s Y_s^T)_{11}}.$$  

Notice that the term proportional to $S_{k1}$ comes from the interference of the tree-level diagram with diagram (d).
For \( x \gg 1 \), the contributions of the self-energy diagrams become negligibly small, and thus the asymmetry can be approximately written as [17],

\[
\varepsilon_1 \simeq -\frac{3M_{R_1}}{16\pi v^2} \text{Im}[(Y_{\nu}^\ast m_\nu Y_{\nu}^\dag)_{11}],
\]

and thus bounded as

\[
|\varepsilon_1| < \frac{3}{16\pi} \frac{M_{R_1}}{v^2} (m_{\nu_3} - m_{\nu_1}),
\]

where \( m_\nu (i = 1 - 3) \) stands for mass eigenvalues of three light neutrinos. For hierarchical neutrino mass spectrum, we can identify \( m_{\nu_3} \simeq \sqrt{\Delta m_{\text{atm}}^2} \) and then it is required that \( M_{R_1} \geq 2 \times 10^9 \) GeV to achieve successful leptogenesis.

On the other hand, for \( x \simeq 1 \), the vertex contribution to \( \varepsilon_1 \) is much smaller than the contribution of the self-energy diagrams and the asymmetry \( \varepsilon \) is resonantly enhanced. To see how much the new contribution can be important in this case, for simplicity, we consider a particular situation where \( M_{R_1} \simeq M_{R_2} < M_{R_3} \), so that the effect of \( N_3 \) is negligibly small. In this case, the asymmetry can be written as

\[
\varepsilon_1 \simeq -\frac{1}{16\pi} \left[ \frac{M_{R_2}}{v^2} \text{Im}[(Y_{\nu}^\ast m_\nu Y_{\nu}^\dag)_{11}] \right. \\
\left. + \sum_{k \neq 1} \frac{\text{Im}[(Y_{\nu} Y_{\nu}^\dag)_{kk} (Y_{s} Y_{s}^\dag)_{kk}]}{(Y_{\nu} Y_{\nu}^\dag + Y_{s} Y_{s}^\dag)_{11}} \right] R,
\]

where \( R \) is a resonance factor defined by \( R \equiv |M_{R_1}|/(|M_{R_2}| - |M_{R_1}|) \). For successful leptogenesis, the size of the denominator of \( \varepsilon_1 \) should be constrained by the out-of-equilibrium condition, \( \Gamma_{N_1} < H|_{T=M_{R_1}} \) with the Hubble expansion rate \( H \), from which the corresponding upper bound on the couplings \( (Y_s)_{1i} \) reads

\[
\sqrt{\sum_i |(Y_s)_{1i}|^2} < 3 \times 10^{-4} \sqrt{M_{R_1}/10^9(\text{GeV})}.
\]

As shown in Eqs. (14, 15), the first term of Eq. (16) is bounded as \( (M_{R_2}/16\pi v^2)\sqrt{\Delta m_{\text{atm}}^2} R \). If the first term of Eq. (16) dominates over the second one, \( R \sim 10^{6-7} \) is required to achieve TeV scale leptogenesis, which implies severe fine-tuning. However, since the size of \( (Y_s)_{2i} \) is not constrained by the out-of-equilibrium condition, large value of \( (Y_s)_{2i} \) is allowed for which the second term of Eq. (16) can dominate over the first one and thus the size of \( \varepsilon_1 \) can be enhanced. For example, if we assume that \( (Y_{\nu})_{2i} \) is aligned to \( (Y_{s})_{2i} \), i.e. \( (Y_s)_{2i} = \kappa (Y_{\nu})_{2i} \) with constant \( \kappa \), the upper limit of the second term of Eq. (16) is given in terms of \( \kappa \) by \( |\kappa|^2 M_{R_2} \sqrt{\Delta m_{\text{atm}}^2} R/16\pi v^2 \), and then we can achieve the successful low scale leptogenesis by
taking rather large value of $\kappa$ instead of imposing very tiny mass splitting between $M_{R_1}$ and $M_{R_2}$.

As one can estimate, the successful leptogenesis can be achieved for $M_{R_1} \sim$ a few TeV, provided that $\kappa = (Y_s)_{2i}/(Y_\nu)_{2i}^* \sim 10^3$ and $M_{R_2}^2/M_{R_1}^2 \sim 10$. We emphasize that such a requirement for the hierarchy between $Y_\nu$ and $Y_s$ is much less severe than the required fine-tuning of the mass splitting between two heavy Majorana neutrinos to achieve the successful leptogenesis at low scale.

The generated B-L asymmetry is given by $Y_{B-L}^{SM} = -\eta \varepsilon_1 Y_{N_1}^{eq}$, where $Y_{N_1}^{eq}$ is the number density of the right-handed heavy neutrino at $T \gg M_{R_1}$ in thermal equilibrium given by $Y_{N_1}^{eq} \simeq \frac{45}{\pi^4 g_{*} k_B} \zeta(3) \frac{3}{4}$ with Boltzmann constant $k_B$ and the effective number of degree of freedom $g_*$. The efficient factor $\eta$ can be computed through a set of coupled Boltzmann equations which take into account processes that create or washout the asymmetry. To a good approximation the efficiency factor depends on the effective neutrino mass $\tilde{m}_1$ given in the presence of the new Yukawa interactions with the coupling $Y_s$ by

$$\tilde{m}_1 = \frac{(Y_\nu Y_\nu^\dagger + Y_s Y_s^\dagger)_{11}}{M_{R_1}} v^2. \quad (18)$$

In our model, the new process of type $S \Psi \rightarrow l \phi$ will contribute to wash-out of the produced B-L asymmetry. The process occurs through virtual $N_{2,3}$ exchanges and the rate is proportional to $M_{R_1} |Y_s Y_\nu^*/M_{R_{2,3}}|^2$. Effect of the wash-out can be easily estimated from the fact that it looks similar to the case of the typical seesaw model if $M_{R_1}$ is replaced with $M_{R_1} (Y_s/Y_\nu)^2$. It turns out that the wash-out factor for $(Y_s)_{1i} \sim (Y_\nu)_{1i}$, $(Y_s)_{2i}/(Y_\nu)_{2i} \sim 10^3$ and $M_{R_1} \sim 10^4$ GeV is similar to the case of the typical seesaw model with $M_{R_1} \sim 10^4$ GeV and $\tilde{m}_1 \simeq 10^{-3}$ eV, and is estimated so that $\varepsilon_1 \sim 10^{-6}$ can be enough to explain the baryon asymmetry of the universe provided that the initial number of the lightest right-handed neutrino is thermal \[18\]. Detailed numerical calculation of the wash-out effect is beyond the scope of this letter and will be presented elsewhere.

Now, let us examine how both light neutrino masses of order 0.01 to 0.1 eV and sterile neutrino masses of order 100 MeV can be simultaneously realized in our scenario without being in conflict with the constraints on the mixing $\theta_s$. Although the absolute values of three neutrinos are unknown, the largest and the next largest neutrino masses are expected to be of order of $\sqrt{\Delta m^2_{atm}} \simeq 0.05$ eV and $\sqrt{\Delta m^2_{sol}} \simeq 0.01$ eV, provided that the mass spectrum of neutrinos is hierarchical. There is also a bound on neutrino mass coming from the WMAP observation, which is $m_\nu \lesssim 0.23$ eV assuming three almost degenerate neutrinos. Thus, it
is interesting to see how the neutrino mass of order of 0.01 ∼ 0.1 eV can be obtained in our scenario. We note that the low scale seesaw can be achieved by taking \( m_D \) to be 1-100 MeV which are of order of the first and second charged lepton masses. To begin with, for our numerical analysis, we take \( \sin^2 \theta_s \simeq 10^{-9} \), which is allowed by the constraints for the mass of sterile neutrinos \( m_s \sim \) a few 100 MeV as described before. Then, we determine the values of the parameters so that the sterile neutrino mass \( m_s \) and the lightest heavy Majorana neutrino mass \( M_{R_1} \) are reached to be about 250 GeV and 1 TeV, respectively.

**Case A** : For \( M^2 > 4\mu M_R \), \( \sin^2 \theta_s \simeq (m_D/M)^2 \) and it follows from Eq. (1) that \( m_{\nu_i} \simeq 0.5 \sin^2 \theta_s \mu_i \). Then the value of the light neutrino mass is determined to be \( m_{\nu_i} \simeq 0.01 \) (0.1) eV for \( \mu_i \simeq 20 \) (200) MeV. For the given value of \( \sin^2 \theta_s \), the size of \( M_i \) should be \( m_D \times \sqrt{10^9} \), and thus the value of the lightest singlet neutrino mass \( M_1 \) is determined to be 30 GeV for \( m_{D_1} \sim 1 \) MeV which is order of electron mass. In this case, the lightest neutrino mass \( m_{\nu_1} \) depends on the size of \( \mu_1 \) which should be much less than 20 MeV in the case of hierarchical light neutrino mass spectrum. We also see from Eq. (6) that \( m_s \simeq 250 \) MeV can be realized by taking \( M_{R_1} \simeq 1 \) TeV. As presented in Eq. (8), the size of \( \mu \) does not strongly affect the value of \( m_s \) in this case.

As shown before, the successful leptogenesis could be achieved for \( M^2 > 4\mu M_R \), \( \sin^2 \theta_s \simeq (m_D/M)^2 \) and it follows from Eq. (1) that \( m_{\nu_i} \simeq 0.5 \sin^2 \theta_s \mu_i \). Then the value of the light neutrino mass is determined to be \( m_{\nu_i} \simeq 0.01 \) (0.1) eV for \( \mu_i \simeq 20 \) (200) MeV. For the given value of \( \sin^2 \theta_s \), the size of \( M_i \) should be \( m_D \times \sqrt{10^9} \), and thus the value of the lightest singlet neutrino mass \( M_1 \) is determined to be 30 GeV for \( m_{D_1} \sim 1 \) MeV which is order of electron mass. In this case, the lightest neutrino mass \( m_{\nu_1} \) depends on the size of \( \mu_1 \) which should be much less than 20 MeV in the case of hierarchical light neutrino mass spectrum. We also see from Eq. (6) that \( m_s \simeq 250 \) MeV can be realized by taking \( M_{R_1} \simeq 1 \) TeV. As presented in Eq. (8), the size of \( \mu \) does not strongly affect the value of \( m_s \) in this case.

**Case B** : For \( M^2 = 4\mu M_R \), \( \tan 2\theta_s \simeq 2 \sin \theta_s \simeq m_D/M \). Then, for \( \sin^2 \theta_s \simeq 10^{-9} \), the value of \( (m_D/M)^2 \) becomes \( 4 \times 10^{-9} \), from which the value of the light neutrino mass is determined to be \( m_{\nu_i} \simeq 0.01 \) (0.1) eV for \( \mu_i \simeq 5 \) (50) MeV. The sizes of \( M_i \) is given by \( 1.6 \times 10^4 m_{D_i} \). In this case, since \( m_{s_i} \simeq 2\mu_i \) as presented in Eq. (9), one can achieve \( m_s \simeq 100 \) MeV which corresponds to \( m_{\nu_i} \simeq 0.1 \), whereas the mass of \( m_s \) corresponding to \( m_{\nu_1} \simeq 0.01 \) is determined to be at most of order 10 MeV. Thus, the case of hierarchical light neutrino spectrum disfavors 100 MeV sterile neutrinos for the first and second generations.

In particular, the size of \( M_R \) is given by \( M^2/(4\mu) \simeq 6 \times 10^7 m_{D_i}^2/\mu \simeq 0.12 m_{D_i}^2/m_{\nu} \), where we used Eqs. (10) to obtain the last relation. For example, if we take \( m_D \simeq 1 \) MeV and \( \nu \simeq 0.1 \) eV, we obtain \( M_R \simeq 1.2 \) TeV. Therefore, it is rather easy to achieve the low scale leptogenesis in consistent with neutrino data as well as 100 MeV sterile neutrino in the case
of the quasi-degenerate light neutrino mass spectrum of order 0.1 eV.

**Case C**: For $4\mu M_R > M^2$, $\tan 2\theta_s \approx 2 \sin \theta_2 \approx m_D M/(2\mu M_R)$. Combining this with Eq. (4), we obtain

$$\sin \theta_s = \frac{m^3_D}{8m_\nu M M_R}. \tag{19}$$

Then, the size of $MM_R$ should be $4 \times 10^5$ $(4 \times 10^{11})$ GeV$^2$ for $\sin^2 \theta_s \approx 10^{-9}$ and $m_D = 1 \ (100)$ MeV. In this case, we have to know the size of $M_i$ as well as $\mu_i$ in order to determine the light neutrino masses. Recall that the value of $m_s$ strongly depends on that of $\mu$ as long as $4\mu M_R >> M^2$.

For smaller values of $\theta_s$, we note that larger value of $\mu$ is demanded so as to achieve the required value of the light neutrino mass. In what follows, we present our numerical results for $\sin^2 \theta_s = 10^{-10}$. Then, $m_{\nu_i} \approx 0.5 \times 10^{-10} \mu_i$, and the values of the light neutrino mass are determined to be $m_{\nu_i(=2,3)} \approx 0.01 \ (0.1)$ eV for $\mu_i \approx 200 \ (2 \times 10^3)$ MeV in Case A and for $\mu_i \approx 50 \ (500)$ MeV in Case B. The size of $M_i$ should be $10^5 \cdot m_{D_i}$ in Case A and $5 \times 10^4 \cdot m_{D_i}$ in Case B. Therefore, taking $m_{D_i} = 1$ MeV , we obtain $M_i \approx 100$ GeV for Case A and $M_i \approx 50$ GeV for Case B. In this case, we can obtain from Eq. (6) that $m_{s_1} \approx 250$ MeV for $M_{R_1} \approx 10$ TeV in Case A, which is desirable for the low scale leptogenesis. For Case B, 100 MeV sterile neutrinos can only be achieved for the second generation. But, we see that the sterile neutrino masses become GeV scale for the cases of $\mu_i = 2 \times 10^3 (500)$ MeV in Case A (B).

In summary, we have considered a variant of seesaw mechanism by introducing extra singlet neutrinos and investigated how the low scale leptogenesis is realized without imposing the tiny mass splitting between two heavy Majorana neutrinos, which makes us avoid the so-called gravitino problem when our scenario is supersymmetrized. We have shown that the introduction of the new singlet fermion leads to a new contribution to lepton asymmetry and it can be enhanced for certain range of parameters. We have also examined how both the light neutrino mass spectrum and relatively light sterile neutrinos of order a few 100 MeV can be achieved without being in conflict with the constraints on the mixing between the active and sterile neutrinos. One of the noticeable features of our seesaw model is that the effects of the heavy right-handed neutrinos are cancelled in seesaw for the tiny neutrino masses, whereas they play an important role in generating light sterile neutrinos.
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