Large lepton asymmetry for small baryon asymmetry and warm dark matter

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We propose a resonant leptogenesis scenario in a $U(1)_{B-L}$ gauge extension of the standard model to generate large lepton asymmetries for cosmological baryon asymmetry and dark matter. After the $B - L$ number is spontaneously broken, the inflaton can pick up a small vacuum expectation value for the mass splits of three pairs of quasidegenerately heavy Majorana neutrinos and the masses of three sterile neutrinos. With thermal mass effects of sphalerons, the observed small baryon asymmetry can be converted from large lepton asymmetries of individual flavors, although total lepton asymmetry is assumed zero. The mixing between sterile and active neutrinos is elegantly suppressed by the heavy Majorana neutrinos. Before the active neutrinos start their strong flavor conversions, the sterile neutrinos as warm dark matter can be produced by resonant active-sterile neutrino oscillations to reconcile x-ray and Lyman-α bounds. Small neutrino masses are naturally realized by seesaw contributions from the heavy Majorana neutrinos and the sterile neutrinos.

I. INTRODUCTION

Sterile neutrinos can provide the dark matter relic density through their oscillations with active neutrinos [1–8]. Because of the mixing with the active neutrinos, the sterile neutrinos can decay at tree level and loop orders. In particular, the decays of a sterile neutrino into an active neutrino and a photon at one-loop order will produce a narrow line in the x-ray background [6]. The x ray constraints put an upper bound on the sterile neutrino mass [9]. Furthermore, the analysis on the Lyman-α data shows a low bound on the sterile neutrino mass [10]. The x ray bound will be in conflict with the Lyman-α bound if the sterile neutrino is produced by the nonresonant oscillations between the active and sterile neutrinos [10]. In the presence of a large neutrino asymmetry, the resonant active-sterile oscillations [2] can reconcile the two bounds [11]. Such a large lepton asymmetry seems inconsistent with the small baryon asymmetry of the Universe because the lepton and baryon asymmetries are usually enforced to be at a same order by sphalerons [12]. This problem can be evaded in three ways: (1) the lepton asymmetry is generated below the electroweak scale; (2) the sphaleron transition does not work; (3) one type of lepton asymmetry is canceled by an opposite lepton asymmetry of other flavors. These possibilities have been proposed and studied in many works [13–17].

On the other hand, observations of solar, atmospheric, reactor, and accelerator neutrino oscillations have established the massive and mixing neutrinos [18]. The cosmological bound shows the neutrino masses should be in the sub-eV range [19]. The smallness of neutrino masses can be naturally explained in the seesaw [20] extension of the standard model (SM). The seesaw essence is to make the neutrino masses tiny via a suppressed ratio of the electroweak scale over a high scale. Most popular seesaw schemes need lepton number violation as the neutrinos are assumed to be Majorana particles. In the seesaw context, the cosmological baryon asymmetry can be understood via leptogenesis [21–25], where a lepton asymmetry is first produced and then is partially converted to a baryon asymmetry through the sphaleron [12] transition. In particular, the so-called resonant leptogenesis [23,24] models [26–28] with quasidegenerately decaying particles can induce a CP asymmetry of the order of the unit. This allows the production of a large lepton asymmetry if we do not take the observed small baryon asymmetry into account.

In this paper, we show that by the resonant leptogenesis the large neutrino asymmetries of individual flavors can be generated for the production of the sterile neutrino dark matter. The total lepton asymmetry is assumed zero so that the baryon asymmetry can arrive at a correct value through the sphaleron processes with thermal mass effects [29]. We demonstrate this possibility in a $U(1)_{B-L}$ gauge extension of the SM. There are two singlet and a doublet Higgs scalars with lepton numbers besides the SM one. One Higgs singlet drives the spontaneous symmetry breaking of the $U(1)_{B-L}$. The other one is responsible for the chaotic inflation [30]. After the $B - L$ number is spontaneously broken, the inflaton can pick up a seesaw suppressed vacuum expectation value (VEV) due to the small ratio of the $B - L$ breaking scale over its heavy mass. In the fermion sector, we introduce three types of SM singlet fermions including the usual right-handed neutrinos. Through the large VEV for the $B - L$ symmetry breaking, the right-handed neutrinos can mix with one type of additional singlet, which gets small Majorana masses from the VEV of the inflaton. So, we can naturally have three pairs of quasidegenerately heavy Majorana neutrinos for the

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resonant leptogenesis. The other singlet fermions obtain small masses through their Yukawa couplings with the inflaton. They can play the role of the sterile neutrinos as their mixing with the active neutrinos is seesaw suppressed. The resonant active-sterile oscillation with neutrinos is induced by a positive neutrino asymmetry, while that with antineutrinos is induced by a negative neutrino asymmetry. Although the total neutrino asymmetry is vanishing, the resonant sterile-active oscillations are still available as they occur much earlier than the beginning of the strong flavor conversions of the active neutrinos. As for the neutrino masses, they are generated in a seesaw scenario, where the heavy Majorana neutrinos and the light sterile neutrinos give a dominant and a negligible contribution, respectively.

II. THE MODEL

The field content of our model is summarized in Table I, where (q_L, d_R, u_R) and (ψ_L, l_R), respectively, are the SM quarks and leptons, φ is the SM Higgs doublet, ν_R denotes usual right-handed neutrinos, ξ_R and ζ_R are additional right-handed singlet fermions, χ and σ are two Higgs singlets, and η is a new Higgs doublet. Note the new fermions ν_R, ξ_R, and ζ_R all have three generations so that the model can be free of gauge anomaly. The full Lagrangian should be $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ invariant. For simplicity we only write down the part relevant for our discussions,

$$
\begin{aligned}
\mathcal{L} \supset &-y_\nu \bar{\psi}_L \varphi \nu_R - y_\xi \bar{\psi}_L \eta \xi_R - f\chi \bar{\nu}_R^c \xi_R - \tfrac{1}{2}h_\xi \sigma^* \bar{\xi}_R^c \xi_R \\
&- \tfrac{1}{2}h_\zeta \sigma \bar{\zeta}_R^c \zeta_R - \mu \bar{\zeta}_R^c \xi_R - \kappa \sigma \chi \varphi^\dagger \eta - \rho \sigma^* \chi^2 \\
&+ \text{H.c.} - m_\sigma^2 |\sigma|^2 - \lambda |\sigma|^2.
\end{aligned}
\tag{1}
$$

TABLE I. Quantum number assignments. Here (q_L, d_R, u_R) and (ψ_L, l_R), respectively, are the SM quarks and leptons, φ is the SM Higgs doublet, ν_R denotes usual right-handed neutrinos, ξ_R and ζ_R are additional right-handed singlet fermions, χ and σ are two Higgs singlets, η is a new Higgs doublet. All of the fermions have three families.

	$SU(3)_c \times SU(2)_L \times U(1)_Y$	$U(1)_{B-L}$
q_L	$(\mathbf{3}, \mathbf{2}, +\tfrac{1}{6})$	$+\tfrac{1}{3}$
d_R	$(\mathbf{3}, \mathbf{1}, -\tfrac{1}{3})$	$+\tfrac{1}{3}$
u_R	$(\mathbf{3}, \mathbf{1}, +\tfrac{2}{3})$	$+\tfrac{1}{3}$
ψ_L	$(\mathbf{1}, \mathbf{2}, -\tfrac{1}{2})$	-1
l_R	$(\mathbf{1}, \mathbf{1}, -1)$	-1
φ	$(\mathbf{1}, \mathbf{2}, -\tfrac{1}{2})$	0
ν_R	$(\mathbf{1}, \mathbf{1}, 0)$	-1
ξ_R	$(\mathbf{1}, \mathbf{1}, 0)$	$+\tfrac{1}{2}$
ζ_R	$(\mathbf{1}, \mathbf{1}, 0)$	$-\tfrac{1}{2}$
χ	$(\mathbf{1}, \mathbf{1}, 0)$	$+\tfrac{1}{2}$
σ	$(\mathbf{1}, \mathbf{1}, 0)$	$+1$
η	$(\mathbf{1}, \mathbf{2}, -\tfrac{1}{2})$	$-\tfrac{3}{2}$

Here the Yukawa couplings h_ξ and h_ζ are symmetric. In addition, the scalar parameters κ and ρ are rotated to be real without loss of generality. Furthermore, the mass term m_σ^2 is positive.

III. HEAVY MAJORANA NEUTRINOS

The $U(1)_{B-L}$ gauge symmetry will be spontaneously broken after the Higgs singlet χ develops its VEV,

$$
\langle \chi \rangle = v_\chi.
\tag{2}
$$

The other Higgs singlet σ will then pick up a small VEV,

$$
\langle \sigma \rangle = v_\sigma \simeq -\frac{\rho v_\chi^2}{m_\sigma^2},
\tag{3}
$$

like the type-II seesaw [31] and its variation [32]. At this stage, we can derive the following mass terms from Eq. (1),

$$
\begin{aligned}
\mathcal{L} \supset &-f v_\chi \bar{\nu}_R^c \xi_R - \tfrac{1}{2}h_\xi v_\sigma \bar{\xi}_R^c \xi_R - \tfrac{1}{2}h_\zeta v_\sigma \bar{\zeta}_R^c \zeta_R \\
&- \mu \bar{\zeta}_R^c \xi_R + \text{H.c.}
\end{aligned}
\tag{4}
$$

The influence of ζ_R on ν_R and ξ_R is negligible for $f v_\chi \gg \mu$. We thus focus on the mass terms only involving ν_R and ξ_R. For illustration, we choose the base where f is diagonal and real, i.e., $f = \text{diag}\{f_1, f_2, f_3\}$. For $f v_\chi \gg h_\xi v_\sigma$, we can perform the following rotation,

$$
\nu_{R_i} \simeq \frac{1}{\sqrt{2}}(N_{R_i}^+ - iN_{R_i}^-),
\tag{5a}
$$

$$
\xi_{R_i} \simeq \frac{1}{\sqrt{2}}(N_{R_i}^+ + iN_{R_i}^-),
\tag{5b}
$$

to obtain the diagonal masses,

$$
\mathcal{L} \supset -\tfrac{1}{2}\bar{N}_R^{+c} m_{N^+} N_R^+ - \tfrac{1}{2}\bar{N}_R^{-c} m_{N^-} N_R^- + \text{H.c.}
\tag{6}
$$

with

$$
m_{N_i^+} \simeq f_i v_\chi + \tfrac{1}{2}h_{\xi_{ii}} v_\sigma,
\tag{7a}
$$

$$
m_{N_i^-} \simeq f_i v_\chi - \tfrac{1}{2}h_{\xi_{ii}} v_\sigma.
\tag{7b}
$$

It is then convenient to define the Majorana fermions,

$$
N_i^+ = N_{R_i}^+ + N_{R_i}^{+c},
\tag{8a}
$$

$$
N_i^- = N_{R_i}^- + N_{R_i}^{-c}.
\tag{8b}
$$

Clearly, N_i^+ and N_i^- are quasidegenerate as their mass split is much smaller than their masses.

We now derive the Yukawa couplings of the heavy Majorana fermions $N^\pm$. For convenience, we first define

$$
\varphi = c\phi_1 + s\phi_2, \qquad \eta = -s\phi_1 + c\phi_2
\tag{9}
$$

with

$$
c \equiv \cos\vartheta, \qquad s \equiv \sin\vartheta, \qquad \vartheta = \frac{1}{2}\arctan\frac{2\kappa v_\chi v_\sigma}{\mu_\varphi^2 - \mu_\eta^2}.
\tag{10}
$$

Here μ_φ^2 and μ_η^2 are the mass terms of φ and η, respectively. The mass terms of ϕ_1 and ϕ_2 would be

$$\mu_{\phi_1}^2 = \frac{1}{2}\{\mu_\varphi^2 + \mu_\eta^2 - [(\mu_\varphi^2 - \mu_\eta^2)^2 + 4\kappa^2 v_\chi^2 v_\sigma^2]^{1/2}\}, \quad (11a)$$

$$\mu_{\phi_2}^2 = \frac{1}{2}\{\mu_\varphi^2 + \mu_\eta^2 + [(\mu_\varphi^2 - \mu_\eta^2)^2 + 4\kappa^2 v_\chi^2 v_\sigma^2]^{1/2}\}. \quad (11b)$$

At least one of μ_φ^2 and μ_η^2 should be negative to guarantee the electroweak symmetry breaking, i.e.,

$$v = \sqrt{v_\varphi^2 + v_\eta^2} \simeq 174 \text{ GeV} \quad \text{with}$$

$$v_\varphi = \langle\varphi\rangle > \frac{m_t}{\sqrt{4\pi}} \simeq \frac{171.2 \text{ GeV}}{\sqrt{4\pi}} \simeq 48 \text{ GeV},$$

$$v_\eta = \langle\eta\rangle. \quad (12)$$

For example, we can take $\mu_\varphi^2 < 0 < \mu_\eta^2$ and then obtain $\mu_{\phi_1}^2 < \mu_\varphi^2 < 0 < \mu_\eta^2 < \mu_{\phi_2}^2$. By inserting Eqs. (5), (8), and (9) into the first and second terms of Eq. (1), we eventually obtain

$$\mathcal{L} \supset -(y_\pm^a)_{\alpha i} \bar{\psi}_{L_\alpha} \phi_a N_i^\pm + \text{H.c.} \quad (13)$$

with

$$y_+^1 = \frac{1}{\sqrt{2}}(cy_\nu - sy_\xi), \quad y_-^1 = -\frac{i}{\sqrt{2}}(cy_\nu + sy_\xi), \quad (14a)$$

$$y_+^2 = \frac{1}{\sqrt{2}}(sy_\nu + cy_\xi), \quad y_-^2 = -\frac{i}{\sqrt{2}}(sy_\nu - cy_\xi). \quad (14b)$$

Clearly, the heavy Majorana fermions $N^\pm$ play the same role with the heavy Majorana neutrinos in the usual seesaw model. We thus refer to $N^\pm$ as the heavy Majorana neutrinos.

IV. RESONANT LEPTOGENESIS

From Eq. (7), it is straightforward to see the heavy Majorana neutrinos N_i^+ and N_i^- have a very small mass split compared to their masses. This means the Yukawa interaction (13) is probably ready for the resonant leptogenesis [23,24] if other conditions are satisfied. Because of the special texture of the Yukawa couplings (14), the decays of $N_i^\pm$ cannot generate a nonzero lepton asymmetry if the two Higgs doublets $\phi_{1,2}$ both appear in the final states. This could be easily understood in the base with (ν_{R_i}, ξ_{R_i}) and (φ, η). Since we have ignored the small Majorana mass term $h_\xi v_\sigma$ for giving the rotation (5) and then the Yukawa couplings (14), both ν_{R_i} and ξ_{R_i} only have one decay channel, i.e., $\nu_{R_i} \to \psi_L + \varphi^*$, $\xi_{R_i} \to \psi_L + \eta^*$. In the presence of the $\varphi - \eta$ mixing (9), we further have the decay channels, $\nu_{R_i} \to \psi_L + \eta^*$, $\xi_{R_i} \to \psi_L + \varphi^*$. Clearly, the decays into $\psi_L \varphi^*$, $\psi_L^c \varphi$ and those into $\psi_L \eta^*$, $\psi_L^c \eta$ will produce an equal but opposite lepton asymmetry stored in the lepton doublets if the CP is not conserved. For a successful leptogenesis, we thus need one of $\phi_{1,2}$ to be heavier than the lightest pair of $N_{1,2,3}^\pm$. For example, we choose ϕ_2 to be heavier than $N_1^\pm$. The other

heavy Majorana neutrinos $N_{2,3}^\pm$, which are assumed much heavier than $N_1^\pm$, have flexibilities to be heavier or lighter than ϕ_2. Therefore a final lepton asymmetry would be produced by the two-body decays of $N_1^\pm$, i.e.,

$$N_1^\pm \to \psi_{L_\alpha} + \phi_1^*, \qquad \psi_{L_\alpha}^c + \phi_1. \quad (15)$$

Following the standard method [24] of the resonant leptogenesis, we can calculate the electron, muon, and tau types of lepton asymmetries from the decays of per $N_1^\pm$,

$$\varepsilon_{N_1^\pm}^{\nu_\alpha} = \varepsilon_{N_1^\pm}^{l_{L_\alpha}} = \frac{1}{2}\varepsilon_{N_1^\pm}^{L_\alpha}$$

$$\simeq \frac{1}{2}\frac{\Gamma(N_1^\pm \to \psi_{L_\alpha} + \phi_1^*) - \Gamma(N_1^\pm \to \psi_{L_\alpha}^c + \phi_1)}{\sum_\alpha[\Gamma(N_1^\pm \to \psi_{L_\alpha} + \phi_1^*) + \Gamma(N_1^\pm \to \psi_{L_\alpha}^c + \phi_1)]}$$

$$\simeq \frac{sc}{16\pi A_{N_1^\pm}}\{[c^2|y_{\nu_{\alpha 1}}|^2 - s^2|y_{\xi_{\alpha 1}}|^2]\text{Im}[(y_\nu^\dagger y_\nu)_{11}]$$

$$+ [c^2(y_\nu^\dagger y_\nu)_{11} - s^2(y_\xi^\dagger y_\xi)_{11}]\text{Im}(y_{\xi_{\alpha 1}}^* y_{\nu_{\alpha 1}})\}$$

$$\times \frac{r_{N_1}}{r_{N_1}^2 + \frac{1}{64\pi^2}A_{N_1^\mp}^2} \quad (16)$$

with

$$A_{N_1^\pm} = \frac{1}{2}\{s^2(y_\xi^\dagger y_\xi)_{11} + c^2(y_\nu^\dagger y_\nu)_{11} \mp sc[(y_\xi^\dagger y_\nu)_{11}$$

$$+ (y_\nu^\dagger y_\xi)_{11}]\}$$

$$= \frac{1}{2}\sum_\alpha(s^2|y_{\xi_{\alpha 1}}|^2 + c^2|y_{\nu_{\alpha 1}}|^2$$

$$\mp 2sc|y_{\xi_{\alpha 1}}||y_{\nu_{\alpha 1}}|\cos\delta_{\alpha 1}). \quad (17)$$

Here the parameter r_{N_i} describes the mass split between N_i^+ and N_i^-,

$$r_{N_i} = \frac{m_{N_i^+}^2 - m_{N_i^-}^2}{m_{N_i^+} m_{N_i^-}} = \frac{2h_{\xi_{ii}} v_\sigma}{f_i v_\chi}. \quad (18)$$

The baryon and lepton asymmetries are determined by the $B - L$ asymmetry in the presence of sphalerons [12]. In the present model, we have [29]

$$B = \frac{28}{79}\left(B - \sum_\alpha L_\alpha\right) = -\frac{28}{79}\sum_\alpha L_\alpha = -\frac{56}{79}\sum_\alpha L_{\nu_\alpha}. \quad (19)$$

The masses and interactions will give corrections to the above formula [29],

$$\Delta B = -A\frac{6}{13\pi^2}\sum_\alpha \frac{\bar{m}_\alpha^2(T)}{T^2}\left(L_\alpha - \frac{1}{3}B\right)$$

$$= -A\frac{6}{13\pi^2}\sum_\alpha \frac{\bar{m}_\alpha^2(T)}{T^2}L_\alpha = -A\frac{12}{13\pi^2}\sum_\alpha \frac{\bar{m}_\alpha^2(T)}{T^2}L_{\nu_\alpha}. \quad (20)$$

Here the coefficient $A \simeq 1$ [29]. In the case that the sphaleron is still active after a weakly first-order electroweak phase transition [33], one finds [29]

$$\frac{\bar{m}_\alpha^2(T)}{T^2} = \frac{1}{6}f_\alpha^2 + \frac{1}{3}f_\alpha^2\left[\frac{\upsilon(T)}{T}\right]^2 \leq \frac{1}{2}f_\alpha^2 \tag{21}$$

for $B - \Sigma_\alpha L_\alpha = 0$. Here f_α denotes the Yukawa couplings of the electron, muon, and tau to the SM Higgs.

With the thermal mass effects of the sphaleron processes, it is possible to generate a small observed baryon asymmetry from large lepton asymmetries of individual flavors [16]. For this purpose, we take

$$\text{Im}[(y_\xi^\dagger y_\nu)_{11}] = \sum_\alpha |y_{\xi_{\alpha 1}}||y_{\nu_{\alpha 1}}|\sin\delta_{\alpha 1} = 0 \tag{22}$$

to give a zero total lepton asymmetry. Under this assumption, the *CP* asymmetry (16) can be simplified by

$$\varepsilon_{N_1^\pm}^{\nu_\alpha} = \varepsilon_{N_1^\pm}^{l_{L_\alpha}} = \frac{1}{2}\varepsilon_{N_1^\pm}^{L_\alpha}$$

$$= \frac{sc}{16\pi A_{N_1^\pm}}[c^2(y_\nu^\dagger y_\nu)_{11} - s^2(y_\xi^\dagger y_\xi)_{11}]$$

$$\times \sin\delta_{\alpha 1}|y_{\xi_{\alpha 1}}||y_{\nu_{\alpha 1}}|\frac{r_{N_1}}{r_{N_1}^2 + \frac{1}{64\pi^2}A_{N_1^\mp}^2}. \tag{23}$$

In the weak washout region where the out-of-equilibrium condition is described by the quantity

$$K_{N_1^\pm} = \frac{\Gamma_{N_1^\pm}}{2H(T)}\bigg|_{T = m_{N_1^\pm}} < 1 \tag{24}$$

with the decay width

$$\Gamma_{N_1^\pm} \simeq \sum_\alpha[\Gamma(N_1^\pm \to \psi_{L_\alpha} + \phi_1^*) + \Gamma(N_1^\pm \to \psi_{L_\alpha}^c + \phi_1)]$$

$$= \frac{1}{8\pi}(y_\pm^{1\dagger}y_\pm^1)_{11}m_{N_1^\pm} = \frac{1}{8\pi}A_{N_1^\pm}m_{N_1^\pm} \tag{25}$$

and the Hubble constant

$$H(T) = \left(\frac{8\pi^3 g_*}{90}\right)^{1/2}\frac{T^2}{M_{\text{Pl}}}, \tag{26}$$

the final neutrino asymmetry can be approximately given by [34]

$$\eta_{\nu_\alpha} = \frac{n_{\nu_\alpha} - n_{\bar\nu_\alpha}}{s} \simeq \frac{\varepsilon_{N_1^\pm}^{\nu_\alpha}}{g_*}. \tag{27}$$

Here $g_* \simeq 112$ is the relativistic degrees of freedom (the SM fields plus $\zeta_{R_{i,2,3}}$), while $M_{\text{Pl}} \simeq 1.22 \times 10^{19}$ GeV is the Planck mass. The final baryon asymmetry then should be

$$\eta_B = \frac{n_B}{s} \simeq -\frac{6}{13\pi^2}\sum_\alpha f_\alpha^2 \eta_{\nu_\alpha}. \tag{28}$$

Note the lepton number violating processes mediated by the heavier $N_{2,3}^\pm$ should be decoupled before the leptogenesis epoch $T = M_{N_1^\pm}$ to give the solution (27).

The neutrino asymmetry η_{ν_α} is related to the neutrino chemical potential μ_{ν_α} by

$$\eta_{\nu_\alpha} = \frac{15}{4\pi^2 g_{*S}(T)}\xi_{\nu_\alpha} + \mathcal{O}(\xi_{\nu_\alpha}^3)$$

$$\simeq \frac{15}{64\pi^2}\xi_{\nu_\alpha} \quad \text{with} \quad \xi_{\nu_\alpha} = \frac{\mu_{\nu_\alpha}}{T_{\nu_\alpha}}. \tag{29}$$

Here we have taken $g_{*S}(T) = 16$ (photon, three neutrinos, electron, positron plus three sterile neutrinos). The electron neutrino asymmetry is tightly constrained by the primordial big-bang nucleosynthesis [35],

$$\eta_{\nu_e} \in (-0.9, 1.7) \times 10^{-3} \quad \text{for } \xi_{\nu_e} \in (-0.04, 0.07). \tag{30}$$

The above bound also applies to the muon and tau neutrino asymmetries because the neutrino oscillations will begin at 10 MeV to achieve strong flavor conversions before the primordial big-bang nucleosynthesis [36]. On the other hand, the five-year observations of the WMAP Collaboration precisely measured the baryon asymmetry as [19]

$$\eta_B = \frac{1}{7.04} \times (6.225 \pm 0.170) \times 10^{-10}$$

$$= (0.884 \pm 0.024) \times 10^{-10}. \tag{31}$$

V. STERILE AND ACTIVE NEUTRINOS

During the evolution of the Universe, the electroweak symmetry breaking will happen when the Higgs doublet

$$H = \frac{\upsilon_\varphi}{\upsilon}\varphi + \frac{\upsilon_\eta}{\upsilon}\eta \tag{32}$$

develops its VEV. The mass terms (4) then should be extended to

$$\mathcal{L} \supset -y_\nu \upsilon_\varphi \bar\nu_L \nu_R - y_\xi \upsilon_\eta \bar\nu_L \xi_R - f\upsilon_\chi \bar\nu_R^c \xi_R$$

$$- \tfrac{1}{2}h_\xi \upsilon_\sigma \bar\xi_R^c \xi_R - \tfrac{1}{2}h_\zeta \upsilon_\sigma \bar\zeta_R^c \zeta_R - \mu \bar\zeta_R^c \xi_R + \text{H.c.} \tag{33}$$

For convenience, we rewrite the above mass terms to be

$$\mathcal{L} \supset -\tfrac{1}{2}(\bar\nu_L, \bar\zeta_R^c, \bar\nu_R^c, \bar\xi_R^c)\mathcal{M}(\nu_L^c, \zeta_R, \nu_R, \xi_R)^T + \text{H.c.}, \tag{34}$$

where the mass matrix $\mathcal{M}$ is defined by

$$\mathcal{M} = \begin{pmatrix} 0 & 0 & y_\nu \upsilon_\varphi & y_\xi \upsilon_\eta \\ 0 & h_\zeta \upsilon_\sigma & 0 & \mu \\ y_\nu^T \upsilon_\varphi & 0 & 0 & f\upsilon_\chi \\ y_\xi^T \upsilon_\eta & \mu^T & f^T \upsilon_\chi & h_\xi \upsilon_\sigma \end{pmatrix}. \tag{35}$$

For $f\upsilon_\chi \gg y_\nu \upsilon_\varphi, y_\xi \upsilon_\eta, h_\zeta \upsilon_\sigma$, and μ, we can make use of the seesaw formula to diagonalize the above mass matrix $\mathcal{M}$ in two blocks,

$$\mathcal{L} \supset -\tfrac{1}{2}(\bar\nu_R^c, \bar\xi_R^c)\mathfrak{M}(\nu_R, \xi_R)^T$$

$$- \tfrac{1}{2}(\bar\nu_L, \bar\zeta_R^c)\mathfrak{m}(\nu_L^c, \zeta_R)^T + \text{H.c.}, \tag{36}$$

where the mass matrices $\mathfrak{M}$ and $\mathfrak{m}$ are given by

$$\mathfrak{M} = \begin{pmatrix} 0 & f v_\chi \\ f^T v_\chi & h_\xi v_\sigma \end{pmatrix}, \qquad \mathfrak{m} = \begin{pmatrix} y_\nu \frac{1}{f^T} h_\xi \frac{1}{f} y_\nu^T \frac{v_\sigma v_\varphi^2}{v_\chi^2} - \left(y_\nu \frac{1}{f^T} y_\xi^T + y_\xi \frac{1}{f} y_\nu^T \right) \frac{v_\varphi v_\eta}{v_\chi} & -y_\nu \frac{1}{f^T} \mu^T \frac{v_\varphi}{v_\chi} \\ -\mu \frac{1}{f} y_\nu^T \frac{v_\varphi}{v_\chi} & h_\zeta v_\sigma \end{pmatrix}. \tag{37}$$

Clearly, the heavy mass matrix $\mathfrak{M}$ will give us the quasi-degenerately heavy Majorana neutrinos $N^\pm$. As for the light mass matrix $\mathfrak{m}$, it can also accommodate the seesaw if its diagonal element $h_\zeta v_\sigma$ is much bigger than other elements. In this seesaw scenario, the neutrino masses should be

$$\mathcal{L} \supset -\tfrac{1}{2}\bar\nu_L m_\nu \nu_L^c + \text{H.c.} = -\tfrac{1}{2}\bar\nu_L(m_\nu^N + m_\nu^S)\nu_L^c + \text{H.c.} \tag{38}$$

with

$$m_\nu^N = y_\nu \frac{1}{f^T} h_\xi \frac{1}{f} y_\nu^T \frac{v_\sigma v_\varphi^2}{v_\chi^2}$$
$$- \left(y_\nu \frac{1}{f^T} y_\xi^T + y_\xi \frac{1}{f} y_\nu^T \right) \frac{v_\varphi v_\eta}{v_\chi}, \tag{39a}$$

$$m_\nu^S = y_\nu \frac{1}{f^T} \mu^T \frac{1}{h_\zeta} \mu \frac{1}{f} y_\nu^T \frac{v_\varphi^2}{v_\sigma v_\chi^2}. \tag{39b}$$

The above neutrino mass matrices are expected to explain the neutrino oscillation experiments [18],

$$\Delta m_{21}^2 = 7.65^{+0.23}_{-0.20} \times 10^{-5} \text{ eV}^2, \qquad \sin\theta_{12}^2 = 0.304^{+0.022}_{-0.016},$$
$$|\Delta m_{31}^2| = 2.4^{+0.12}_{-0.11} \times 10^{-3} \text{ eV}^2, \qquad \sin\theta_{23}^2 = 0.50^{+0.07}_{-0.06},$$
$$\sin\theta_{13}^2 = 0.01^{+0.016}_{-0.011}. \tag{40}$$

On the other hand, ζ_R plays the role of the sterile neutrinos, i.e.,

$$\mathcal{L} \supset -\tfrac{1}{2}h_\zeta v_\sigma \bar\zeta_R^c \zeta_R + \text{H.c.} = -\tfrac{1}{2}m_S \bar S S, \tag{41}$$

where we have rotated $h_\zeta = \text{diag}\{h_{\zeta_1}, h_{\zeta_2}, h_{\zeta_3}\}$ and then defined

$$S_i = \zeta_{R_i} + \zeta_{R_i}^c \quad \text{with} \quad m_{S_i} = h_{\zeta_i} v_\sigma. \tag{42}$$

The active-sterile mixing angle is

$$\theta_{\alpha i}^2 = \frac{\left| y_{\nu_{\alpha j}} \frac{1}{f_j} \mu_{ij} \frac{v_\varphi}{v_\chi} \right|^2}{m_{S_i}^2} \quad \text{and} \quad \sin^2 2\theta_i = 4\sum_\alpha \theta_{\alpha i}^2. \tag{43}$$

The sterile neutrinos can be produced through the active-sterile neutrino oscillations. Specifically, the $\nu_\alpha \to S_i$ or $\bar\nu_\alpha \to S_i$ oscillation is determined by the sterile neutrino mass m_{S_i}, the active-sterile mixing angle $\theta_{\alpha i}$, and the neutrino asymmetry η_{ν_α}. In order to compare with other works, we define the following function [4,37],

$$c_{\alpha\alpha} = \sqrt{2}G_F \left[2\eta_{\nu_\alpha} + \sum_{\beta \neq \alpha} \eta_{\nu_\beta} + (1 + 2\sin^2\theta_W)\eta_{l_{L_\alpha}} \right.$$
$$\left. - (1 - 2\sin^2\theta_W)\sum_{\beta \neq \alpha} \eta_{l_{L_\beta}} + 2\sin^2\theta_W \sum_\beta \eta_{l_{R_\beta}} \right]$$
$$= 3\sqrt{2}G_F \eta_{\nu_\alpha}. \tag{44}$$

Here we have taken $\eta_{\nu_\alpha} = \eta_{l_\alpha}$, $\sum_\alpha \eta_{\nu_\alpha} = 0$, and $\eta_{l_R} = 0$ into account. A positive neutrino asymmetry $\eta_{\nu_\alpha} > 0$ can induce a Mikheyev-Smirnov-Wolfenstein [38] resonant behavior in the neutrino oscillation $\nu_\alpha \to S_i$, whereas a negative one $\eta_{\nu_\alpha} < 0$ can enhance the antineutrino oscillation $\bar\nu_\alpha \to S_i$ [2,37]. It has been pointed out that a large neutrino asymmetry can reconcile the contradiction between the x ray and Lyman-α bounds. For example, in the work [39,40] where the lepton asymmetry is flavor blind, i.e., $\eta_{\nu_\alpha} = \eta_{l_{L_\alpha}} = \eta_{l_{R_\alpha}} = \eta_{\nu_e}$, the authors show with the parameter choice,

$$m_{S_i} \simeq 8 \text{ keV}, \qquad \sin^2 2\theta_i \simeq 10^{-12}, \qquad \eta_{\nu_e} = 7 \times 10^{-5}, \tag{45}$$

the resonant $\nu_\alpha \to S_i$ oscillation can produce adequate S_i for the dark matter relic density. We should keep in mind that the resonant active-sterile conversion happens at a temperature of the order of 100 MeV for the sterile neutrino mass m_{S_i} being a few keV [3,4]. This means the electron neutrino asymmetry η_{ν_e} and the opposite muon and tau neutrino asymmetries $\eta_{\nu_{\mu,\tau}}$ can survive for the resonant enhancement either in the neutrino oscillations or in the antineutrino oscillations since the active neutrino oscillations driven by Δm_{31}^2 and Δm_{21}^2 begin at lower temperatures $T \sim 10$ MeV and $T \sim 3$ MeV, respectively.

VI. PARAMETER CHOICE

We now take a reasonable choice of parameters to give the observed small cosmological baryon asymmetry from large lepton asymmetries for the dark matter production. First, the heavy Majorana neutrinos $N_i^\pm$ couple to the gauge boson Z_{B-L} associated with the $U(1)_{B-L}$ symmetry. For $m_{N_1^\pm} = \mathcal{O}(10^5 \text{ GeV})$, the Higgs singlet χ should have a VEV bigger than $\mathcal{O}(10^8 \text{ GeV})$ to guarantee the departure from equilibrium of $N_1^\pm$ at $T \simeq M_{N_1^\pm}$, unless the corresponding gauge coupling g_{B-L} is fine-tuned very small [41]. So we take

$$v_\chi = \mathcal{O}(10^8 \text{ GeV}) \tag{46}$$

and then

$$m_{N_1^\pm} = 10^5 \text{ GeV} \quad \text{for } f_1 = \mathcal{O}(10^{-3}),$$
$$m_{N_{2,3}^\pm} = 10^6 \text{ GeV} \quad \text{for } f_{2,3} = \mathcal{O}(10^{-2}). \tag{47}$$

Second, the Higgs singlet σ is expected to realize the nice picture of the chaotic inflation [30]. This suggests [42]

$$m_\sigma = \mathcal{O}(10^{13} \text{ GeV}), \qquad \lambda = \mathcal{O}(10^{-13}). \tag{48}$$

We then conveniently set

$$\rho = v_\chi \tag{49}$$

in Eq. (3) to induce a small VEV,

$$v_\sigma = \mathcal{O}(10 \text{ MeV}). \tag{50}$$

In consequence, the mass splits of the heavy Majorana neutrinos $N_i^\pm$ can be determined by

$$r_{N_i} = \mathcal{O}(10^{-12}) \quad \text{for } h_{\xi_{11}} = \mathcal{O}(10^{-5}),$$
$$h_{\xi_{22,33}} = \mathcal{O}(10^{-4}). \tag{51}$$

The sterile neutrinos S_i also obtain small masses:

$$m_{s_i} = \mathcal{O}(1\text{–}100 \text{ keV}) \quad \text{for } h_{\zeta_i} = \mathcal{O}(10^{-4}\text{–}10^{-2}). \tag{52}$$

Third, we consider

$$-[\mathcal{O}(10^2 \text{ GeV})] = \mu_\varphi^2 < 0 < \mu_\eta^2 = [\mathcal{O}(10^5 \text{ GeV})]^2, \tag{53}$$

to derive the mixing angle,

$$\vartheta \simeq -\frac{\kappa v_\chi v_\sigma}{\mu_\eta^2} = 10^{-4} \quad \text{for } \kappa = \mathcal{O}(1), \tag{54}$$

and then determine the VEVs:

$$v_\varphi \simeq 174 \text{ GeV}, \qquad v_\eta \simeq -\frac{\kappa v_\chi v_\sigma v_\varphi}{\mu_\eta^2} \simeq v_\varphi \vartheta. \tag{55}$$

We now consider the following sample of the Yukawa couplings,

$$|y_{\nu_{e1}}| = 1.15 \times 10^{-6}, \qquad |y_{\xi_{e1}}| = 10^{-3}, \qquad \sin\delta_{e1} = 1,$$

$$|y_{\nu_{\mu1}}| = 1.15 \times 10^{-7}, \qquad |y_{\xi_{\mu1}}| = \frac{10^{-2}}{2},$$

$$\sin\delta_{\mu1} = -1, \qquad |y_{\nu_{\tau1}}| = 1.15 \times 10^{-7},$$

$$|y_{\xi_{\tau1}}| = \frac{10^{-2}}{2}, \qquad \sin\delta_{\mu1} = -1, \tag{56}$$

to fulfill the assumption (22). We then read

$$K_{N_1^+} = K_{N_1^-} = 0.135. \tag{57}$$

By further fixing $r_{N_1} = 10^{-12}$, we can obtain the CP asymmetry,

$$\varepsilon_{N_1^+}^{\nu_e} = \varepsilon_{N_1^-}^{\nu_e} = 2.03 \times 10^{-3}. \tag{58}$$

In consequence, the large lepton asymmetries should be

$$\eta_{\nu_e} = 3.63 \times 10^{-5} = \frac{6.99 \times 10^{-5}}{1 + 4\sin^2\theta_W}, \tag{59}$$

which is in agreement with the observations (30). Accordingly a desired baryon asymmetry is induced by

$$\eta_B \simeq \frac{6}{13\pi^2} \frac{f_\mu^2 + xf_\tau^2}{1+x} \eta_{\nu_e} = 0.888 \times 10^{-10} \tag{60}$$

$$\text{for } x = \frac{|y_{\nu_{\tau1}}||y_{\xi_{\tau1}}|\sin\delta_{\tau1}}{|y_{\nu_{\mu1}}||y_{\xi_{\mu1}}|\sin\delta_{\mu1}} = 1.$$

We further take

$$y_{\nu_{\alpha2}} = y_{\nu_{\alpha3}} = 100y_{\nu_{\alpha1}}, \qquad \mu_{1j} = 20 \text{ keV}, \tag{61}$$

and then perform

$$4\theta_{e1}^2 = 10^{-12}, \qquad 4\theta_{\mu1}^2 = 4\theta_{\tau1}^2 = 10^{-14}. \tag{62}$$

So the $\nu_e \to S_1$ oscillation is dominant. Compared with the fitting by [39,40], we see the sterile neutrino S_1 with the mass

$$m_{S_1} \simeq 8 \text{ keV} \tag{63}$$

can be a good candidate for the dark matter. The other sterile neutrinos $S_{2,3}$ are also at the keV scale. They could leave a significant relic density if their mixing with the active neutrinos $\theta_{\alpha i}^2 (i = 2, 3)$ are comparable with $\theta_{\alpha1}^2$. Alternatively, their relic density could be negligible if $\theta_{\alpha i}^2 (i = 2, 3)$ are much smaller than $\theta_{\alpha1}^2$. This could be achieved by taking appropriate $\mu_{ij}(i = 2, 3)$ in Eq. (43).

For the above parameter choice, the seesaw contribution from the sterile neutrinos to the neutrino masses is of the order of $\theta_{\alpha i}^2 m_{S_i} \lesssim \mathcal{O}(10^{-9} \text{ eV})$, which is too small to explain the neutrino oscillation data. Furthermore, the contribution from the heavy Majorana neutrinos $N_1^\pm$ is also too small [$\sim \mathcal{O}(10^{-5} \text{ eV})$]. Fortunately, the other heavy Majorana neutrinos $N_{2,3}^\pm$ can give a desired contribution. Specifically, we can take the Yukawa couplings $y_{\xi_{\alpha i}} \lesssim \mathcal{O}(1)(i = 2, 3)$ in the neutrino mass matrix which is dominated by the second term of Eq. (39a). We also check the lepton number violating processes mediated by $N_{2,3}^\pm$ have been decoupled before the leptogenesis epoch since their reaction rate [43],

$$\Gamma_A = \frac{1}{\pi^3} \frac{\text{Tr}(m_\nu^\dagger m_\nu)}{v^4} T^3 = \frac{1}{\pi^3} \frac{\sum_i m_i^2}{v^4} T^3, \tag{64}$$

is smaller than the Hubble constant at the temperature $T = M_{N_1^\pm}$.

VII. SUMMARY

We have shown the resonant leptogenesis in the seesaw context for the observed small baryon asymmetry can allow large neutrino asymmetries of individual flavors for the production of the sterile neutrino dark matter due to the

thermal mass effects of the sphaleron processes. In our model, the inflaton can nicely obtain a small VEV as a result of the seesaw suppressed ratio of the $B - L$ breaking scale over its heavy mass. The quasidegenerate mass spectrum of the heavy Majorana neutrinos for the resonant leptogenesis is naturally given by the scale of the $B - L$ symmetry breaking and the small VEV of the inflaton. The sterile neutrinos obtain small Majorana masses through their Yukawa couplings with the inflaton. Their mixing with the active neutrinos is suppressed by the seesaw mechanism. Thus one or more sterile neutrinos can act as the warm dark matter in the presence of the resonant active-sterile neutrino oscillation. Because of the small active-sterile mixing, the sterile neutrinos only have a negligible contribution to the neutrino masses. Instead, the seesaw contribution from the heavy Majorana neutrinos is responsible for generating the desired neutrino masses.

For an appropriate parameter choice, we may realize the νMSM model [3,14,40] (the SM plus three right-handed neutrinos) with a pair of quasidegenerately heavy Majorana neutrinos and a sterile neutrino in our model with three pairs of quasidegenerately heavy Majorana neutrinos and three sterile neutrinos. In the νMSM model, the processes for generating the lepton asymmetry can keep working below the sphaleron freeze-out temperature at which there is no conversion of the lepton asymmetry to the baryon asymmetry. In this case, we need not resort to the thermal mass effects of the sphaleron processes.

ACKNOWLEDGMENTS

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