

# Leptogenesis and dark matter unified in a non-SUSY model for neutrino masses

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**Abstract.** We propose a unified explanation for the origin of dark matter and baryon number asymmetry on the basis of a non-supersymmetric model for the neutrino masses. Neutrino masses are generated in two distinct ways, that is, a tree-level seesaw mechanism with a single right-handed neutrino, and one-loop radiative effects by a new additional doublet scalar. A spontaneously broken  $U(1)'$  brings about a  $Z_2$  symmetry which restricts couplings of this new scalar and controls the neutrino masses. It also guarantees the stability of a CDM candidate. We examine two possible candidates for the CDM. We also show that the decay of a heavy right-handed neutrino related to the seesaw mechanism can generate baryon number asymmetry through leptogenesis.

## 1 Introduction

Neutrino masses [1–11], cold dark matter (CDM) [12, 13], and baryon number asymmetry in the universe [14] suggest that the standard model (SM) should be extended. Both neutrino masses and baryon number asymmetry are well known to be explainable in a unified way through the leptogenesis scenario in the framework of the seesaw mechanism [15]. Extensive studies have been done on this subject for recent years [16–21]. On the other hand, supersymmetry is known to play a crucial role for the explanation of CDM abundance in the universe [22–26], although it has been introduced originally to solve the hierarchy problem. Supersymmetric models have good candidates for CDM such as the lightest superparticle (LSP) as long as  $R$ -parity is conserved. The neutralino LSP has been extensively studied as a CDM candidate in the supersymmetric SM (MSSM) and its singlet extensions [27–40]. If we try to explain simultaneously both leptogenesis and the CDM abundance in supersymmetric models, we have a difficulty. The out-of-equilibrium decay of thermal heavy neutrinos can generate a sufficient baryon number asymmetry only if the reheating temperature is high enough such as  $T_R > 10^8$  GeV. For such a reheating temperature, however, we are confronted by a serious gravitino problem in supersymmetric models [41, 42]. Various attempts to overcome this difficulty have been done by searching scenarios to enhance the  $CP$  asymmetry and lower the required reheating temperature [43–58].

In these studies, the CDM and the baryon number asymmetry are separately explained based on unrelated physics. Thus, we cannot expect to obtain any hints as to

why the CDM abundance is of similar order as the baryon number asymmetry in the present universe through such studies.<sup>1</sup> Unfortunately, at present, we have no satisfactory supersymmetric models to explain these three experimental pieces of evidence which lead us to extend the SM. In this situation, it may be worthwhile to take a different empirical view point at first and reconsider possible models which can explain this evidence simultaneously on the basis of closely related physics [67, 68]. As the next step, the hierarchy problem may be considered in the framework where such models are embedded.

Recently, it has been suggested that neutrino masses and the CDM abundance may be related in some kind of non-supersymmetric models for neutrino masses. In such models neutrino masses are generated through one-loop radiative effects which are induced by new scalar fields [69]. A certain  $Z_2$  symmetry prohibiting large neutrino masses can also guarantee the stability of a CDM candidate like  $R$ -parity in supersymmetric models [70–74]. The baryon number asymmetry has also been discussed in this model [75, 76]. In the same type of model there is also a suggestion that the hierarchy problem can be improved by considering a heavy Higgs scalar [77]. Since these models have a rather simple structure at weak scale regions, it might give us some useful hints for physics beyond the SM, if they can explain the above mentioned experimental evidence consistently.

In this paper, we consider the possibility that the baryon number asymmetry is closely related to the origin of both neutrino masses and CDM abundance. We show

<sup>1</sup> There are several works to relate the CDM abundance to the baryon number asymmetry. For such attempts, see [59–66] for example.

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that the ordinary leptogenesis based on heavy neutrino decay can be embedded consistently in the model for neutrino masses proposed in [74]. As we discuss below, this is closely related to an extension of [74] such that (1) an additional  $N$  with zero charge under  $U(1)'$  is introduced and (2) the dimension five term in the scalar potential has a complex coupling  $\lambda_6$ . The paper also includes new contributions added to [74] in so far that (1) both  $N_3$  and  $\eta_0$  are studied as dark matter candidates and (2) the constraints due to neutrino oscillation data are taken into account in a more extended way than that in [74].

The remaining parts are organized as follows. In Sect. 2, we address features of the model and discuss a parameter space consistent with neutrino oscillation data. In Sect. 3 we study the relation between leptogenesis and the CDM abundance in the model. We examine two possible CDM candidates taking account of the neutrino oscillation data and the conditions required by the leptogenesis. We shall find that the model can give a unified picture for the explanation of the neutrino masses, the CDM abundance, and the baryon number asymmetry. In Sect. 4, we summarize the paper with comments on the signatures of the model expected at LHC.

## 2 A model for neutrino masses

The present study is based on the model proposed in [74]. Ingredients of the model and  $U(1)'$  charge assignments for these are given in Table 1. We suppose that  $U(1)'$  is leptophobic.<sup>2</sup> The extension to general  $U(1)'$  is straightforward. The fermions listed in Table 1 are assumed to be left-handed. We note that three singlet fermions  $N_{1,2,3}$  are necessary for present purposes. Although only two of them are ordered to generate appropriate masses and mixing in the neutrino sector, an additional one is necessary for the leptogenesis. The invariant Lagrangian relevant to the neutrino masses can be expressed by

$$\begin{aligned} \mathcal{L}_m = & \sum_{\alpha=e,\mu,\tau} (h_{\alpha 1} L_\alpha H \bar{N}_1 + h_{\alpha 2} L_\alpha H \bar{N}_2 + h_{\alpha 3} L_\alpha \eta \bar{N}_3) \\ & + \frac{1}{2} M_1 \bar{N}_1^2 + \frac{1}{2} M_2 \bar{N}_2^2 + \frac{1}{2} \lambda \phi \bar{N}_3^2 + \text{h.c.} \end{aligned} \quad (1)$$

Yukawa couplings for charged leptons are assumed to be diagonalized already. The most general scalar potential invariant under  $SU(2) \times U(1) \times U(1)'$  gauge symmetry up to dimension five is given by

$$\begin{aligned} V = & \frac{1}{2} \lambda_1 (H^\dagger H)^2 + \frac{1}{2} \lambda_2 (\eta^\dagger \eta)^2 + \frac{1}{2} \lambda_3 (\phi^\dagger \phi)^2 \\ & + \lambda_4 (H^\dagger H)(\eta^\dagger \eta) + \lambda_5 (H^\dagger \eta)(\eta^\dagger H) \\ & + \frac{1}{2M_*} [\lambda_6 \phi (\eta^\dagger H)^2 + \text{h.c.}] \\ & + (m_H^2 + \lambda_7 \phi^\dagger \phi) H^\dagger H + (m_\eta^2 + \lambda_8 \phi^\dagger \phi) \eta^\dagger \eta + m_\phi^2 \phi^\dagger \phi, \end{aligned} \quad (2)$$

<sup>2</sup> We need to introduce some fields to cancel the gauge anomalies. However, it can be done without affecting the following study. We present such an example in the appendix.

**Table 1.** Field contents and their charges.  $Z_2$  is the residual symmetry of  $U(1)'$

	$Q_\alpha$	$\bar{U}_\alpha$	$\bar{D}_\alpha$	$L_\alpha$	$\bar{E}_\alpha$	$\bar{N}_{1,2}$	$\bar{N}_3$	$H$	$\eta$	$\phi$
$U(1)'$	$2q$	$-2q$	$-2q$	$0$	$0$	$0$	$q$	$0$	$-q$	$-2q$
$Z_2$	$+1$	$+1$	$+1$	$+1$	$+1$	$+1$	$-1$	$+1$	$-1$	$+1$

where the couplings  $\lambda_i$  are real except for  $\lambda_6$ . The phase of  $\lambda_6$  can induce a physical one which is found to be a Majorana phase in the neutrino mass matrix. A nonrenormalizable  $\lambda_6$  term and bare mass terms for  $N_{1,2}$  are added, which will be shown to play crucial roles in the present scenario. They are supposed to be effective terms generated through some dynamics at intermediate scales. We assume that  $M_* \simeq M_1 \ll M_2$  and only  $N_1$  and  $N_3$  are related to light neutrino masses and mixings.

The model includes two  $SU(2)$  doublet scalars  $H$  and  $\eta$ .  $H$  plays the role of the ordinary doublet Higgs scalar in the SM, but  $\eta$  is assumed to obtain no VEV. A singlet scalar  $\phi$  is also assumed to have a real VEV at suitable scales, which breaks  $U(1)'$  down to  $Z_2$ . The  $Z_2$  charge for each field can be found in Table 1. The VEV of  $\phi$  gives masses for  $N_3$  and  $Z'$ :

$$M_{N_3} = \lambda \langle \phi \rangle, \quad M_{Z'} = 2\sqrt{2}g'q \langle \phi \rangle, \quad (3)$$

where  $\lambda$  is assumed to be real. Since  $M_{Z'}$  is bounded from below by the  $Z'$  phenomenology,  $M_{N_3}$  has also lower bounds for fixed values of  $\lambda$ . It also yields an effective coupling constant  $\lambda_6 \langle \phi \rangle / M_*$  in the  $\lambda_6$  term. It may be small enough to make radiative neutrino masses tiny even for  $O(1)$  values of  $\lambda_6$  as long as  $\langle \phi \rangle \ll M_*$  is satisfied. Since the mixing between  $\eta^0$  and  $\eta^{0*}$  is induced through this small coupling, the mass eigenvalues split slightly. The states  $\chi_\pm^0 \equiv \frac{1}{\sqrt{2}} (\eta^0 \pm \eta^{0*})$  have mass eigenvalues such as

$$\begin{aligned} M_{\chi_\pm^0}^2 = & m_\eta^2 + (\lambda_4 + \lambda_5) \langle H^0 \rangle^2 + \lambda_8 \langle \phi \rangle^2 \pm \frac{|\lambda_6| \langle \phi \rangle}{M_*} \langle H^0 \rangle^2 \\ \equiv & M_\eta^2 \pm \frac{|\lambda_6| \langle \phi \rangle}{M_*} \langle H^0 \rangle^2. \end{aligned} \quad (4)$$

The magnitude of the difference of these eigenvalues is constrained by the direct search of the CDM if either of these  $\chi_\pm^0$  is the lightest  $Z_2$  odd field. The mass of the charged states  $\eta^\pm$  is given by

$$M_{\eta^\pm} = m_\eta^2 + \lambda_4 \langle H^0 \rangle^2 + \lambda_8 \langle \phi \rangle^2, \quad (5)$$

and then  $M_{\chi_\pm^0}$  may be much smaller than  $M_{\eta^\pm}$  in case of  $\lambda_5 < 0$ . These points will be discussed in the analysis of the CDM later. Since  $\lambda_6$  is complex in general, the  $CP$  violation may be detected through this  $\eta^0$ - $\eta^{0*}$  mixing. Although this is an interesting feature of the model, we do not discuss this subject further in this paper.

We have two distinct origins for the neutrino masses in this model. One is the ordinary seesaw mass induced by a right-handed neutrino  $N_1$  [78–81]. Another one is the one-loop radiative mass mediated by the exchange of  $\eta^0$

and  $N_3$  [69, 82–84]. Although  $N_2$  also has contributions to the neutrino mass generation through the seesaw mechanism, its effect can be safely neglected compared with these, if  $M_2$  is large enough. However, baryogenesis caused by leptogenesis requires this contribution, since  $N_3$  has no lepton number as discussed below. The radiative neutrino mass generation requires some lepton number violation. We can put them either in  $\mathcal{L}_m$  or  $V$ . If we assume that  $\eta$  and  $N_3$  have the lepton number  $-1$  and  $0$ , respectively, the  $\lambda_6$  term in  $V$  brings about this required lepton number violating effect. We adopt this choice in the following arguments.  $N_{1,2}$  are considered to have lepton number  $+1$ .

The mass matrix for three light neutrinos induced by these origins is summarized as

$$M_\nu = \frac{\langle H^0 \rangle^2}{M_*} \left[ \mu^{(1)} + \frac{\lambda_6}{8\pi^2 \lambda} I \left( \frac{M_{N_3}^2}{M_{\eta_0}^2} \right) \mu^{(3)} \right],$$

$$I(t) = \frac{t}{1-t} \left( 1 + \frac{t \ln t}{1-t} \right), \quad (6)$$

where  $\mu^{(a)}$  is defined by

$$\mu^{(a)} = \begin{pmatrix} h_{ea}^2 & h_{ea} h_{\mu a} & h_{ea} h_{\tau a} \\ h_{ea} h_{\mu a} & h_{\mu a}^2 & h_{\mu a} h_{\tau a} \\ h_{ea} h_{\tau a} & h_{\mu a} h_{\tau a} & h_{\tau a}^2 \end{pmatrix} \quad (a = 1, 3). \quad (7)$$

Both  $h_{\alpha 1}$  and  $h_{\alpha 3}$  are assumed to be real, for simplicity. We note that two terms in  $M_\nu$  have the similar texture although they are characterized by different mass scales. If we impose commutativity between  $\mu^{(1)}$  and  $\mu^{(3)}$ , the condition

$$h_{e1} h_{e3} + h_{\mu 1} h_{\mu 3} + h_{\tau 1} h_{\tau 3} = 0 \quad (8)$$

is to be satisfied. We consider this simple case in the following as an interesting example, since it allows us to study the mass matrix analytically.<sup>3</sup>

We introduce a matrix  $\tilde{U}$  to diagonalize the larger term of  $M_\nu$  at first, which is defined as

$$\tilde{U} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos \theta_2 & \sin \theta_2 \\ 0 & -\sin \theta_2 & \cos \theta_2 \end{pmatrix} \begin{pmatrix} \cos \theta_3 & 0 & \sin \theta_3 \\ 0 & 1 & 0 \\ -\sin \theta_1 & 0 & \cos \theta_3 \end{pmatrix}. \quad (9)$$

Then the matrix  $\mu^{(a)}$  in  $M_\nu$  can be diagonalized as  $\tilde{U}^T \mu^{(a)} \tilde{U}$  if the angles  $\theta_{2,3}$  satisfy

$$\tan \theta_2 = \frac{h_{\mu a}}{h_{\tau a}}, \quad \tan \theta_3 = \frac{h_{ea}}{\sqrt{h_{\mu a}^2 + h_{\tau a}^2}}. \quad (10)$$

Eigenvalues for this matrix are found to be

$$\mu_{\text{diag}}^{(a)} = \text{diag} (0, 0, h_{ea}^2 + h_{\mu a}^2 + h_{\tau a}^2). \quad (11)$$

Another term  $\mu^{(a')}$  is also transformed by  $\tilde{U}$ . However, if the condition (8) is satisfied,  $\mu^{(a')}$  can be diagonalized by

an orthogonal transformation and  $\tilde{U}U_1$  supplemented by the additional transformation

$$U_1 = \begin{pmatrix} \cos \theta_1 & \sin \theta_1 & 0 \\ -\sin \theta_1 & \cos \theta_1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (12)$$

and we have the eigenvalues

$$\mu_{\text{diag}}^{(a')} = \text{diag} (0, h_{ea'}^2 + h_{\mu a'}^2 + h_{\tau a'}^2, 0). \quad (13)$$

Here  $\theta_1$  is defined by

$$\tan \theta_1 = -\frac{\tan \tilde{\theta}_2 \tan \theta_2 + 1}{(\tan \tilde{\theta}_2 - \tan \theta_2) \sin \theta_3}, \quad \tan \tilde{\theta}_2 = \frac{h_{\mu a'}}{h_{\tau a'}}. \quad (14)$$

We note that this  $U_1$  transformation does not affect the diagonalization of  $\mu^{(a)}$ .

If we define the mass eigenvalues by  $U^T M_\nu U = \text{diag}(0, m_2, m_3)$ , where  $m_2 < m_3$  is assumed, they can be written as

$$m_2 = AB \frac{\tan^2 \theta_1 + 1}{\tan^2 \theta_2 + 1} (\tan \tilde{\theta}_2 - \tan \theta_2)^2,$$

$$m_3 = \frac{A}{2} (\tan^2 \theta_2 + 1) (\tan^2 \theta_3 + 1). \quad (15)$$

Here we find that there are two possibilities for the generation of  $m_3$  and  $m_2$ . The first case is realized by taking  $a = 1$  and  $a' = 3$  in the above formulas, and then  $m_3$  is induced by the ordinary seesaw mechanism. In this case  $A$  and  $B$  are defined by

$$(i) \quad A \equiv \frac{2h_{\tau 1}^2 \langle H^0 \rangle^2}{M_*}, \quad B \equiv \frac{|\lambda_6|}{8\pi^2 \lambda} \left( \frac{h_{\tau 3}}{h_{\tau 1}} \right)^2 I \left( \frac{M_{N_3}^2}{M_{\eta_0}^2} \right). \quad (16)$$

The second case is obtained by taking  $a = 3$  and  $a' = 1$ , and then  $m_3$  is determined by the radiative effect. In this case  $A$  and  $B$  are written as

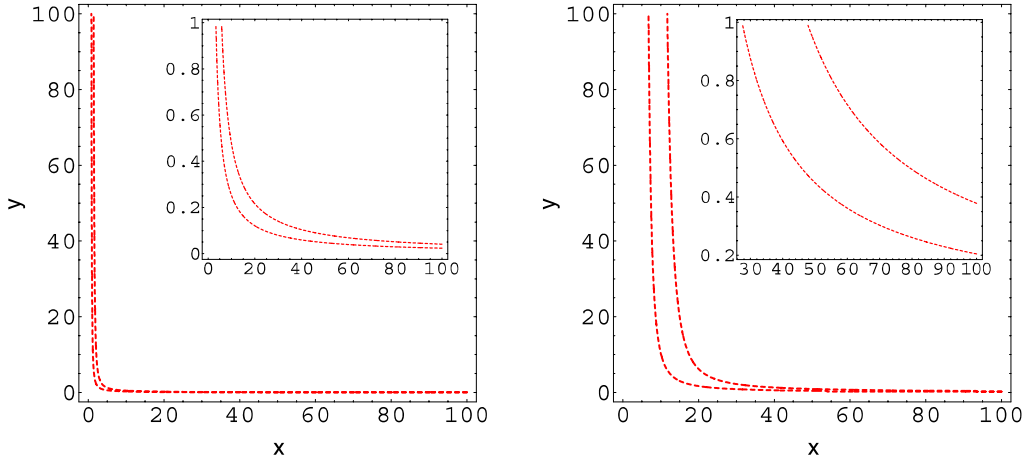
$$(ii) \quad A \equiv \frac{h_{\tau 3}^2 \langle H^0 \rangle^2}{M_*} \frac{|\lambda_6|}{4\pi^2 \lambda} I \left( \frac{M_{N_3}^2}{M_{\eta_0}^2} \right),$$

$$B \equiv \left[ \frac{|\lambda_6|}{4\pi^2 \lambda} I \left( \frac{M_{N_3}^2}{M_{\eta_0}^2} \right) \right]^{-1} \left( \frac{h_{\tau 1}}{h_{\tau 3}} \right)^2. \quad (17)$$

Since only two mass eigenvalues may be considered non-zero in the present setting, neutrino oscillation data require that these mass eigenvalues should satisfy  $m_3 = \sqrt{\Delta m_{\text{atm}}^2}$  and  $m_2 = \sqrt{\Delta m_{\text{sol}}^2}$  [1–11]. Data of the atmospheric neutrino and the K2K experiment require  $\tan \theta_2 = 1$ . We also find that  $\theta_1$  should be taken as  $\theta_{\text{sol}}$ , which is a mixing angle relevant to the solar neutrino. The CHOOZ experiment gives a constraint on  $\theta_3$  such as  $|\sin \theta_3| < 0.22$  [85]. If we use these conditions, the mixing matrix  $U = \tilde{U}U_1$  can be approximately written as

$$U = \begin{pmatrix} \cos \theta_{\text{sol}} & \sin \theta_{\text{sol}} & \frac{\sin \theta_3}{\sqrt{2}} \\ -\frac{\sin \theta_{\text{sol}}}{\sqrt{2}} & \frac{\cos \theta_{\text{sol}}}{\sqrt{2}} & \frac{1}{\sqrt{2}} \\ \frac{\sin \theta_{\text{sol}}}{\sqrt{2}} & -\frac{\cos \theta_{\text{sol}}}{\sqrt{2}} & \frac{1}{\sqrt{2}} \end{pmatrix}. \quad (18)$$

<sup>3</sup> If non-zero eigenvalues are dominated by different origins, respectively, this will be a good approximation to describe such cases.



**Fig. 1.** The region in the  $(x, y)$  plane allowed by the neutrino oscillation data. The cases (i) and (ii) correspond to the *left-handed* and *right-handed* panel, respectively. A figure focusing on the  $0 < y < 1$  region is also displayed in each panel

By imposing the experimental values on  $\tan \theta_{\text{sol}}$ ,  $\sqrt{\Delta m_{\text{atm}}^2}$ ,  $\sqrt{\Delta m_{\text{sol}}^2}$ , and  $\sin \theta_3$ , we can constrain the values of  $A$  and  $B$  [74]. For simplicity, we assume  $\lambda = |\lambda_6|$ .

The condition for  $A$  constrains the Yukawa coupling  $h_{\tau 1}$  by

$$\begin{aligned} \text{(i)} \quad h_{\tau 1} &\simeq 2.9 \times 10^{-4} \left( \frac{M_*}{10^8 \text{ GeV}} \right)^{1/2}, \\ \text{(ii)} \quad 7.9 \times 10^{-5} &\left( \frac{M_*}{10^8 \text{ GeV}} \right)^{1/2} \\ &\lesssim h_{\tau 1} \lesssim 1.3 \times 10^{-4} \left( \frac{M_*}{10^8 \text{ GeV}} \right)^{1/2}. \end{aligned} \quad (19)$$

If we require  $h_{\tau 1}$  and  $h_{\tau 3}$  to be in perturbative regions, we find that both  $M_*$  and  $M_* x^2$  should be less than  $10^{16}$  GeV. Here we introduce two parameters  $x \equiv h_{\tau 3}/h_{\tau 1}$  and  $y \equiv M_{N_3}/M_\eta$ . The condition for  $B$  selects the regions in the  $(x, y)$  plane which are consistent with the neutrino oscillation data. They are shown for both cases (i) and (ii) as the regions sandwiched by the dashed lines in Fig. 1. These figures show that the model may explain the neutrino oscillation data in rather wide parameter regions. In particular, it is useful to note in relation to the CDM that we may have solutions for large values of  $y$  such as  $10^6$  as long as  $x$  stays in the constrained regions (i) 0.55–0.8 and (ii) 3.5–6.5. By using these results obtained from the neutrino oscillation data, we examine the leptogenesis and the CDM abundance in this model in the next section.

### 3 Leptogenesis and CDM abundance

The present model contains several new neutral fields with non-zero lepton number or an odd  $Z_2$  charge. Thus, we have sufficient ingredients with the required properties for both leptogenesis and CDM candidates. Although one might consider that there are several scenarios for these explanations in this model, they seem to be constrained by the neutrino oscillation data.

The lightest neutral field with an odd  $Z_2$  charge may be stable and a CDM candidate, since an even charge is

assigned to each SM content. If  $y < 1$  is satisfied,  $N_3$  can be a CDM candidate. As in the ordinary leptogenesis scenario,  $N_1$ , related to the ordinary seesaw mechanism may be a mother field for leptogenesis. However, since two right-handed neutrinos are necessary to realize the  $CP$  asymmetry, we need to introduce  $N_2$  with the lepton number  $+1$  as mentioned before.

On the other hand, since  $\eta^0$  has both the odd  $Z_2$  charge and the lepton number, it might be considered as the origin of the CDM or the lepton number asymmetry in the case of  $y > 1$ . However, it might be difficult to contribute both of them, since it has the SM gauge interactions. The situation is similar to sneutrinos in the supersymmetric models. Sneutrinos have been rejected to be a CDM candidate through direct detection experiments. This constraint might be escapable in the case of  $\eta^0$ , since there is  $\eta^0 - \eta^{0*}$  mixing due to the  $\lambda_6$  term which generates the mass difference between its components. The model has to satisfy suitable conditions for this mass difference if this possibility is realized. On the other hand, this  $\eta^0$  is too light to be a mother field for the sufficient production of the lepton number asymmetry through the out-of-equilibrium decay, although the  $\eta^0$  sector can bring about almost degenerate mass eigenstates through  $CP$  violating mixing and cause resonant decay. We examine these subjects in detail below.

#### 3.1 Leptogenesis

If we take account of the existence of  $N_2$  which can be neglected in the estimation of the neutrino masses, the leptogenesis is expected to occur through the decay of  $N_1$ . In fact, it is heavy enough for the out-of-equilibrium decay, and it has lepton number violation through a Majorana mass term. By taking account of the well known relation  $B = 28(B - L)/75$  which comes from re-processing of the  $B-L$  asymmetry by sphaleron transitions, the generated baryon number asymmetry is given by

$$\frac{n_B}{s} = -\frac{28}{75} Y_{N_1}^{\text{eq}} \varepsilon \kappa, \quad (20)$$

where  $Y_{N_1}^{\text{eq}} (\equiv n_{N_1}/s)$  is the ratio of the equilibrium number density of  $N_1$  to the entropy density. The  $CP$  asymmetry in

the  $N_1$  decay and the wash-out effect are represented by  $\varepsilon$  and  $\kappa$ , respectively. If temperature is much larger than  $M_1$ , we have  $Y_{N_1}^{\text{eq}} \simeq 0.42/g_*$  by using  $n_{N_1} = (3\zeta(3)/2\pi^2)T^3$  and  $s = (2\pi^2 g_*/45)T^3$ . The relativistic degrees of freedom in this model number  $g_* \simeq 130$ . Thus, the  $CP$  asymmetry  $\varepsilon$  required to produce the present baryon number asymmetry is estimated to be

$$\varepsilon \simeq -7.2 \times 10^{-8} \kappa^{-1}, \quad (21)$$

where we use  $n_B/s \simeq (0.87 \pm 0.04) \times 10^{-10}$ , which is predicted by nucleosynthesis and CMB measurements [14]. The  $CP$  violation in the  $N_1$  decay is induced through interference between the tree and one-loop amplitudes. This induced  $CP$  asymmetry  $\varepsilon$  is estimated to be [16–21]

$$\varepsilon = -\frac{3}{16\pi} \frac{M_1}{M_2} \frac{\text{Im}[(h^\dagger h)_{12}^2]}{|h^\dagger h|_{11}}. \quad (22)$$

Now we estimate  $\varepsilon$  in this model. As discussed in the previous section, there are two ways for the generation of the neutrino masses  $m_3$  and  $m_2$ . The  $CP$  asymmetry  $\varepsilon$  can also have different values for these two cases. For simplicity, we assume  $|h_{\alpha 2}| \simeq |h_{\alpha 1}|$ . This does not affect the estimation of the neutrino masses, because of the assumed setting  $M_* \simeq M_1 \ll M_2$ . In that case, we have

$$|\text{Im}[(h^\dagger h)_{12}^2]| \lesssim 4h_{\tau 1}^4 \simeq \begin{cases} 2.8 \times 10^{-14} \left(\frac{M_*}{10^8 \text{ GeV}}\right)^2 & \text{for (i),} \\ (0.16\text{--}1.1) \times 10^{-14} \left(\frac{M_*}{10^8 \text{ GeV}}\right)^2 & \text{for (ii),} \end{cases} \quad (23)$$

where we apply the results in (19) to this estimation. We use these maximum values for  $\text{Im}[(h^\dagger h)_{12}^2]$  in the formulas of  $\varepsilon$  here.

In case (i), we have the relation  $|h^\dagger h|_{11} \langle H_0 \rangle^2 / M_* \simeq \sqrt{\Delta m_{\text{atm}}^2}$ , and then  $\varepsilon$  can be written as

$$\varepsilon \simeq -9.8 \times 10^{-8} \left(\frac{10^{10} \kappa^{-1} \text{ GeV}}{M_2}\right) \left(\frac{M_*}{10^8 \text{ GeV}}\right)^2 \kappa^{-1}. \quad (24)$$

In case (ii), we note that the seesaw mechanism gives  $m_2$  and the relation  $|h^\dagger h|_{11} \langle H_0 \rangle^2 / M_* \simeq \sqrt{\Delta m_{\text{sol}}^2}$  is satisfied. Thus, we find that  $\varepsilon$  is expressed by

$$\varepsilon = -2.2 \times 10^{-8} \left(\frac{10^{10} \kappa^{-1} \text{ GeV}}{M_2}\right) \left(\frac{M_*}{10^8 \text{ GeV}}\right)^2 \kappa^{-1}. \quad (25)$$

These results show that a sufficient  $CP$  asymmetry can be generated for

$$M_* \simeq \begin{cases} 8.6 \times 10^7 \left(\frac{M_2}{10^{10} \kappa^{-1} \text{ GeV}}\right)^{1/2} \text{ GeV} & \text{for (i),} \\ 1.8 \times 10^8 \left(\frac{M_2}{10^{10} \kappa^{-1} \text{ GeV}}\right)^{1/2} \text{ GeV} & \text{for (ii).} \end{cases} \quad (26)$$

Consistency with the present setting  $M_2 \gg M_*$  can be satisfied for  $M_2 \gtrsim 10^{10} \kappa^{-1} \text{ GeV}$  in both cases, for example. It may be useful to recall that  $\kappa$  is expected to be  $10^{-1}\text{--}10^{-3}$  from a numerical study of the Boltzmann equation. Such an analysis also shows that the leptogenesis is possible only for narrow ranges of  $\tilde{m}_1 = |h^\dagger h|_{11} \langle H_0 \rangle^2 / M_1$  [16–21]. In the present model this  $\tilde{m}_1$  is estimated to be

$$\tilde{m}_1 \simeq \begin{cases} \sqrt{\Delta m_{\text{atm}}^2} \frac{M_*}{M_1} & \text{for (i),} \\ \sqrt{\Delta m_{\text{sol}}^2} \frac{M_*}{M_1} & \text{for (ii).} \end{cases} \quad (27)$$

This suggests that  $M_* \lesssim M_1$  is favored by leptogenesis and it could be consistent in the present settings. The values of  $M_*/M_1$  determine which case between them is more promising. These results show that the out-of-equilibrium decay of  $N_1$  may produce the necessary baryon number asymmetry for intermediate values of  $M_1$  as in the usual cases. As long as we confine ourselves to the non-supersymmetric framework, the model is free from the gravitino problem.

### 3.2 CDM candidates and their abundance

The lightest field with an odd  $Z_2$  charge can be stable since an even charge is assigned to each SM content. If both the mass and the annihilation cross section of such a field have appropriate values, it may be a good CDM candidate as long as it is neutral. As mentioned before, we have two such candidates; that is, the lighter one of  $\chi_{\pm}^0$  (we represent it by  $\chi_L^0$ ) and  $N_3$ .

At first, we consider the  $y < 1$  case in which  $N_3$  is the CDM. Its annihilation is expected to be mediated by both the exchange of  $\eta^0$  and the  $U(1)'$  gauge boson. If their annihilation is mediated only by the former one through Yukawa couplings as in the model discussed in [72, 73], we need fine tuning of the coupling constants to explain both the observed value of the CDM abundance and the constraints coming from lepton flavor violating processes such as  $\mu \rightarrow e\gamma$ . However, in the present case the  $N_3$  annihilation can be dominantly mediated by the  $U(1)'$  gauge interaction, since the Yukawa coupling constants  $h_{\alpha 3}$  can be small enough as estimated in (19). Thus, we may expect that  $N_3$  can cause a satisfactory relic abundance as the CDM in rather wide parameter regions. We also note that the  $U(1)'$  is supposed to be a generation independent gauge symmetry and then the FCNC problem can easily be escaped in this case.

In order to estimate the  $N_3$  abundance, we consider to expand the annihilation cross section for  $N_3 N_3 \rightarrow f\bar{f}$  by the relative velocity  $v$  between the annihilating  $N_3$  as  $\sigma v = a + bv^2$ . The coefficients  $a$  and  $b$  are expressed by

$$a = \sum_f c_f \frac{g^4}{2\pi} Q_{fA}^2 q^2 \frac{m_f^2 \beta}{(s - M_{Z'}^2)^2}, \quad b = \sum_f c_f \frac{g^4}{6\pi} (Q_{fV}^2 + Q_{fA}^2) q^2 \frac{M_{N_3}^2 \beta}{(s - M_{Z'}^2)^2}, \quad (28)$$

where  $\beta = \sqrt{1 - m_f^2/M_{N_3}^2}$  and  $c_f = 3$  for quarks.  $s$  is the center of mass energy of collisions, and  $q$  is the  $U(1)'$  charge of  $N_3$  given in Table 1. The charge of the final state fermion  $f$  is defined by

$$Q_{f_V} = Q_{f_R} + Q_{f_L}, \quad Q_{f_A} = Q_{f_R} - Q_{f_L}. \quad (29)$$

Using these quantities, the present relic abundance of  $N_3$  can be estimated to be [86]

$$\Omega_{N_3} h^2|_0 = \frac{M_{N_3} n_{N_3}}{\rho_{\text{cr}}/h^2} \Big|_0 \simeq \frac{8.76 \times 10^{-11} g_*^{-1/2} x_F}{(a + 3b/x_F) \text{GeV}^2}. \quad (30)$$

where  $g_*$  is the number of degrees of freedom of relativistic fields at the freeze-out temperature  $T_F$  of  $N_3$ . The dimensionless parameter  $x_F = M_{N_3}/T_F$  is determined through the condition

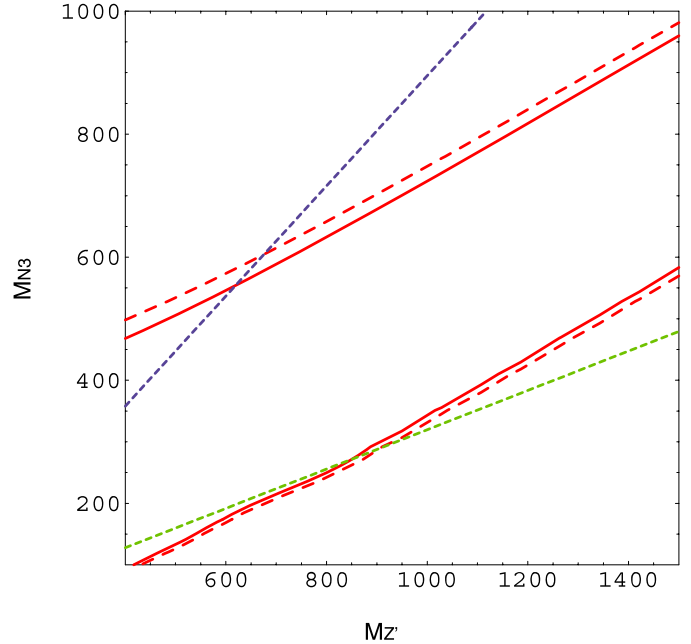
$$x_F = \ln \frac{0.0955 m_{\text{pl}} M_{N_3} (a + 6b/x_F)}{(g_* x_F)^{1/2}}, \quad (31)$$

where  $m_{\text{pl}}$  is the Planck mass. If we fix the  $U(1)'$  charge of the relevant fields and its coupling constant  $g'$ , we can estimate the present  $N_3$  abundance using these formulas. It can be compared with  $\Omega_{N_3} h^2 = 0.1045_{-0.0095}^{+0.0072}$  given by the three year WMAP [87].

We numerically examine the possibility that the CDM abundance is consistently explained in this model. We use the GUT relation  $g' = \sqrt{5/3} g_1$  and  $q = 0.6$  as an example. The regions in the  $(M_{Z'}, M_{N_3})$  plane allowed by the WMAP data are shown in Fig. 2. They appear as two narrow bands sandwiched by both a solid line and a dashed line. The lower bounds of  $M_{Z'}$  come from constraints for  $ZZ'$  mixing and a direct search of  $Z'$ . Since the Higgs field  $H$  is assumed to have no  $U(1)'$  charge, its VEV induces no  $ZZ'$  mixing. Moreover, since it is assumed to be leptophobic, the constraint on  $M_{Z'}$  obtained from its hadronic decay is rather weak. The lower bounds of  $M_{Z'}$  may be  $M_{Z'} \gtrsim 450$  GeV in the present model [88, 89]. Since the masses of  $Z'$  and  $N_3$  are correlated through (3), we can draw a line of  $M_{N_3}$  in the  $(M_{Z'}, M_{N_3})$  plane by fixing a value of  $\lambda$ . In Fig. 2, such lines are represented by the green and blue dotted ones for  $\lambda = 0.25$  and  $0.7$ , respectively. For these  $M_{N_3}$  values as required by the WMAP,  $M_\eta$  is found to take values such as  $\sim 300/y$  and  $\sim 580/y$  GeV for  $\lambda = 0.25$  and  $0.7$ . Using Figs. 1 and 2, we can determine the range of  $x$ , if  $M_\eta$  and then  $y$  is fixed. We find that  $x$  takes very restricted values for the case of  $M_\eta \lesssim 1$  TeV, especially in case (i).

In Fig. 2 we can observe an interesting feature of  $Z'$ . Although we assume that it is leptophobic, it can have a nonhadronic decay model as long as  $2M_{N_3} < M_{Z'}$  is satisfied. Figure 2 shows that this condition is satisfied only at the lower allowed band but not at the upper allowed band. Thus,  $Z'$  may have a nonhadronic decay mode only for  $\lambda \lesssim 0.33$ .

If  $y > 1$  is satisfied, the neutral scalar  $\chi_L^0$  is the CDM. In this case, we can follow the analysis given in [77]. If it is heavier than the  $W^\pm$  boson, it cannot keep the relic abundance required from the WMAP data. The reason is that they can effectively annihilate to the  $W^\pm$  pair through



**Fig. 2.** Regions allowed by the WMAP data in the  $(M_{Z'}, M_{N_3})$  plane. *Green and blue dotted lines* represent  $M_{N_3}$  lines for  $\lambda = 0.25$  and  $0.7$ , respectively

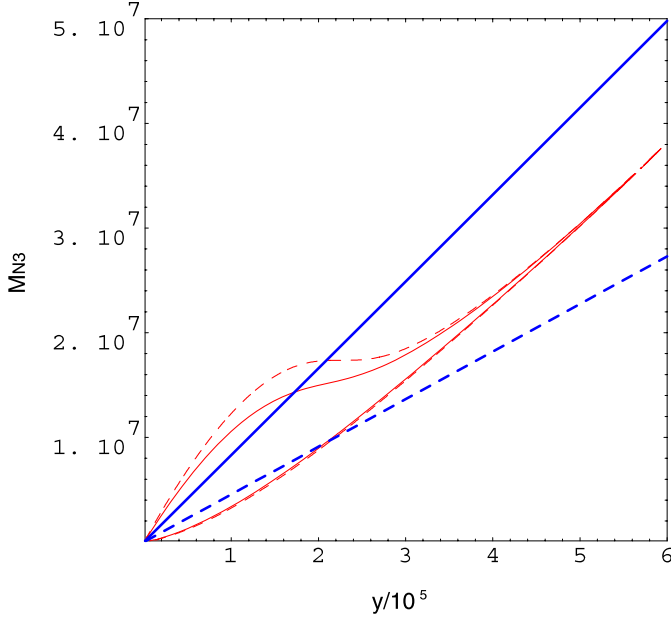
the  $Z^0$  exchange. Thus, since we have no other candidate for the CDM within the present model, we have to assume that the mass of  $\chi_L^0$  should be smaller than 80 GeV. Even if it is lighter than the  $W^\pm$  boson, direct search experiments impose a strong constraint. The difference of the mass eigenvalues of  $\chi_\pm^0$  is estimated to be

$$\begin{aligned} \Delta M &\simeq \frac{|\lambda_6| \langle \phi \rangle \langle H^0 \rangle^2}{M_\eta M_*} \sim \frac{M_{N_3}}{M_\eta M_*} \langle H^0 \rangle^2 \\ &\sim 300y \left( \frac{10^8 \text{ GeV}}{M_*} \right) \text{keV}. \end{aligned} \quad (32)$$

Since the  $\chi_\pm^0$  have a vector-like interaction with the  $Z^0$  boson, its elastic scattering cross section with a nucleon through  $Z^0$  exchange is 8–9 orders of magnitude larger than the existing direct search limits [90]. To forbid  $Z^0$  exchange kinematically,  $\Delta M$  has to be larger than a few 100 keV [91, 92]. Following (32), this constraint can be interpreted as a condition  $y \gtrsim (M_*/10^8 \text{ GeV})$ .

If we impose that the relic  $\chi_L^0$  abundance saturates the values required by the WMAP data, a much stronger constraint can be obtained. This  $\chi_L^0$  abundance is dominantly determined by the  $p$ -wave suppressed coannihilation process  $\chi_+^0 \chi_-^0 \rightarrow Z^* \rightarrow ff$ . In order to realize a suitable relic abundance, we need to decrease this coannihilation rate by requiring the heavier one of  $\chi_\pm^0$  to be thermally suppressed. This requires that  $\Delta M \gtrsim 8\text{--}9$  GeV should be satisfied for  $M_{\chi_L^0} = 60\text{--}73$  GeV [77]. Thus, if we assume that  $\chi_L^0$  is the CDM taking account of these arguments, we have an another condition,  $y \gtrsim M_*/(3000 \text{ GeV})$ . Since the leptogenesis occurs successfully for  $M_* \gtrsim 10^9$  GeV as seen in the previous part,  $y$  should have a larger value than  $2 \times 10^5$ , and then  $M_{N_3}$  should be larger than  $3 \times 10^7$  GeV.





**Fig. 3.** Allowed regions in the  $(y, M_{N_3})$  plane. A red thin dotted line and a red thin solid line corresponds to an upper and lower bound of  $\Omega_{\chi_L^0} h^2$  imposed by the WMAP data. A blue thick solid line represents a line for  $M_{\chi_L^0} = 80$  GeV. A blue thick dotted line represents a boundary for  $M_{\chi_+^0} + M_{\chi_-^0} = m_Z$

We can search favored parameter regions in the present model by numerically estimating the relic abundance of  $\chi_L^0$  in the same way as the  $N_3$  case. In this estimation we need to take account of the above mentioned thermal effect, which modifies the relic density in the  $\Delta M = 0$  case by a factor  $\frac{1}{2} \exp(\Delta M/T_F)$ . In Fig. 3 we plot the allowed regions in the  $(y, M_{N_3})$  plane for the case of  $M_* = 10^9$  GeV, which is a favored value for leptogenesis. In the regions sandwiched by both the dotted and solid thin lines,  $\Omega_{\chi_L^0}$  realizes the three year WMAP data. In the same figure we add two conditions. We plot a line corresponding to  $M_{\chi_L^0} = 80$  GeV by a blue solid thick one. Since we now consider regions below the  $WW$  threshold, the allowed regions are those below this line. The  $Z^0$  width also imposes another condition,  $M_{\chi_+^0} + M_{\chi_-^0} > m_Z$ . The boundary of this condition is plotted by a blue dotted thick line. Regions above this boundary satisfy this condition. As seen from this figure, the favored part in the regions sandwiched by these thick lines gives 40–80 GeV for  $M_{\chi_L^0}$ , which agrees with the results given in [77, 91, 92]. This does not contradict the experimental mass bounds for charged Higgs fields as long as  $\lambda_4$  has suitable negative values. The constraint from  $\mu \rightarrow e\gamma$  can also be satisfied for  $M_*$ , which may keep the Yukawa couplings small enough in (19). For the required large values  $(2-5) \times 10^5$  for  $y$ ,  $|\lambda_6| \langle \phi \rangle \ll M_*$  can still be satisfied, and  $Z'$  becomes very heavy so as to be out of the range reached by the LHC experiments.<sup>4</sup> In

this case,  $x$  is confined to very restricted regions, especially in case (i). In order to realize the favorable values of  $M_{\chi_L^0}$  and  $\Delta M$ , several coupling constants are required to be finely tuned. For example,  $\lambda_8$  should be very small, something like  $O(10^{-5})$ . Although this required parameter tuning might decrease the interests for this case compared with the  $y < 1$  case, it is noticeable that  $\chi_L^0$  can be a CDM candidate consistent with the neutrino oscillation data in this model.

## 4 Summary

We have studied a unified explanation for both the CDM abundance and the baryon number asymmetry in a non-supersymmetric model for the neutrino masses. The model is obtained from the SM by adding a  $U(1)'$  gauge symmetry and several neutral fields. The neutrino masses are generated through both the seesaw mechanism with a single right-handed neutrino and the one-loop radiative effects. The two contributions induce the same texture, which can realize favorable mass eigenvalues and mixing angles. New neutral fields required for this mass generation make the unified explanation for the leptogenesis and the CDM abundance in the universe possible.

Both the neutral fermion  $N_3$  and the neutral scalar  $\eta^0$  are stable due to a  $Z_2$  subgroup which remains as a residual symmetry of the spontaneously broken  $U(1)'$ . Thus, they can be a good CDM candidate. In the  $N_3$  CDM case, since it has the  $U(1)'$  gauge interaction, the annihilation of this CDM candidate is dominantly mediated through this interaction. If this  $U(1)'$  symmetry is broken at a scale suitable for the neutrino mass generation, its estimated relic abundance can explain the WMAP result for the CDM abundance. We examined these points taking account of the neutrino oscillation data. In the  $\eta^0$  CDM case, if it is lighter than the  $W^\pm$  boson and the difference of its mass eigenstates forbid its coannihilation due to  $Z^0$  exchange kinematically, it may keep a suitable relic abundance. We examined the consistency of this picture with the neutrino oscillation data.

Since another introduced neutral fermion,  $N_1$ , is a gauge singlet and is heavy enough, it can follow the out-of-equilibrium decay which produces the baryon number asymmetry through the leptogenesis. We showed the consistency of this scenario with the neutrino oscillation data. Although the required reheating temperature for leptogenesis is similar in value to the one in the ordinary seesaw mechanism, we have no gravitino problem, since we need no supersymmetry to prepare the stable CDM candidates. The present model gives an example in which three of the biggest experimental questions in the SM, that is, neutrino masses, the CDM abundance, and the baryon number asymmetry can be explained through the closely related physics

<sup>4</sup> In the original models [72, 73], the required values of  $\Delta M$  and  $M_{\chi_L^0}$  for the  $\chi_L^0$  CDM can be consistent with the neutrino oscillation data and the FCNC constraint as long as singlet

fermion masses are large enough, and their Yukawa couplings are small as in the present case. Thus, we could not find substantial differences between this model and the original ones in the  $y > 1$  case.

in a non-supersymmetric extension of the SM. In order to solve the hierarchy problem, a supersymmetric extension of the model may be considered along the lines of [93]. We would like to discuss this subject elsewhere.

Finally, we briefly comment on the signatures of the model expected at LHC. The above study fixes the mass spectrum of the relatively light fields in the model. We have  $N_3$ ,  $\eta$  and  $Z'$  as such new fields.  $\eta$  is expected to be produced through the  $W$  fusion in a way similar to the ordinary Higgs field. Since  $\eta$  has Yukawa couplings with leptons only, its components  $\eta^0$  and  $\eta^\pm$  can be distinguished from others such as the Higgs fields in the MSSM through the difference of the decay modes.  $Z'$  couples with the quarks,  $\eta$ , and  $N_3$ . However, its decay shows a different feature, depending on the scheme for the CDM. If the CDM is  $N_3$ , the results shown in Fig. 2 suggest that the decay mode of  $Z'$  is mainly hadronic. It can include nonhadronic ones only for the case of  $\lambda \lesssim 0.33$ , as mentioned before. In such cases, in the  $Z'$  decay  $\ell^+\ell^- +$  a missing energy also is included in the final states depending on the value of  $y$ . On the other hand, if one component of  $\eta^0$  is the CDM, the  $Z'$  always can decay into the  $\eta$  pair, since it is very light. Thus,  $Z'$  has a substantial invisible width. The search of  $Z'$  with such features may be an important check of the model.

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## Appendix

We give an example of a set of fields which cancel the gauge anomalies without affecting the discussion in the text. We consider it appropriate to introduce additional fermions as the left-handed ones:

$$2(\mathbf{3}, 0, -q); \quad 3 \left[ \left( \mathbf{2}, +\frac{1}{2}, -q \right) + \left( \mathbf{2}^*, -\frac{1}{2}, -q \right) \right]; \\ 6[(\mathbf{1}, +1, q) + (\mathbf{1}, -1, q)]; 5(\mathbf{1}, 0, q), \quad (\text{A.1})$$

where representations and charges for  $SU(2) \times U(1)_Y \times U(1)'$  are shown in parentheses. The numbers of fields are also given in front of them. The SM gauge anomalies are canceled by taking account of these fields. Since these fields are vector-like for the SM gauge group, no problem is induced by them as seen against the electroweak precision measurements. Although these fields are  $Z_2$  odd, all of them may be massive through Yukawa couplings with  $\phi$  or  $\phi^*$ . Thus, as long as their Yukawa coupling constants with  $\phi$  or  $\phi^*$  are simply larger than  $\lambda$ ,  $\bar{N}_3$  remains as the lightest  $Z_2$  odd field in the model. Some discrete symmetry such as  $Z_2$  seems to be necessary to forbid the coupling between  $\bar{N}_3$  and singlet fields shown in the last line of (A.1). However, it can be introduced without affecting the scenario. Since no other seeds for the  $U(1)'$  breaking are necessary to make these additional fermions massive, the mass formula for  $m_{Z'}$  does not change, and the discussion on the relic abundance in the text is not affected.

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