

## Dirac leptogenesis in extended nMSSM

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## Dirac leptogenesis in extended nMSSM

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**ABSTRACT:** We show that a version of the nearly Minimal Supersymmetric Standard Model (nMSSM), extended only in the singlet sector to include the additional superfields of right-handed neutrinos and very heavy Dirac particles conserving  $B-L$ , admits a viable scenario for Dirac leptogenesis and naturally small Dirac neutrino masses. The origin of the  $(B-L)$ -conserving high singlet neutrino scale and the desired supersymmetry breaking terms is associated with dynamical supersymmetry breaking in the hidden sector.

**KEYWORDS:** Baryogenesis, Beyond Standard Model, Neutrino Physics, Supersymmetry Phenomenology.

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## 1. Introduction

Leptogenesis, generating a sphaleron-induced baryon asymmetry from lepton asymmetry, can be realized not only with Majorana neutrinos [1] but also with Dirac neutrinos conserving  $B - L$  [2]. Such neutrinos (as well as charged leptons) are first produced in the early universe with equal numbers of left-handed and right-handed states from the decays of very heavy fields. However, the Yukawa couplings of the former are small enough to hide, during the electroweak transition epoch, the right-handed lepton number from sphalerons. This is as opposed to the left-handed lepton number which gets converted to baryon number through sphaleronic transitions.

In the above context, a realistic supersymmetric model for Dirac neutrino masses and leptogenesis was proposed in ref. [3], and studied in detail in refs. [4]. In this model, a sub-eV neutrino mass scale gets generated by a combination of reasonably small Yukawa couplings and a high mass scale of heavy SU(2) doublet superfields. Furthermore, leptogenesis can successfully arise from the decays of such heavy fields. However, there is a serious problem with this kind of a supersymmetric scenario. The same heavy superfields can also decay into right-handed sneutrinos which have masses of the order of hundreds of GeV to a bit more or less than a TeV, as induced by soft SUSY-breaking terms. Since their couplings  $yLHN$ , which are supersymmetric generalizations of Yukawas, are tiny with  $y \sim 10^{-13}$  guaranteeing the non-equilibration of the right-handed neutrinos, they do not decay quickly enough and persist as relics after the electroweak transition. In fact, their number density divided by the entropy density then is expected to be as large as  $10^{-3}$ . If they were stable, their relic density today would be far too much in excess of that expected of Dark Matter (DM) as estimated from cosmological observations. On the other hand, because of their tiny couplings, unstable right-handed sneutrinos can only decay very slowly, thus surviving till long after the decoupling of the Lightest Supersymmetric Particle (LSP) which takes place typically at a temperature of the order of 10 GeV. Their

late decays will then lead to a late non-thermal over-production of the LSPs, once again conflicting with the required DM relic density. Our aim is to present the conditions on supersymmetry breaking parameters in the neutrino sector avoiding this difficulty, and we present here a variant of the Murayama-Pierce model [3] which achieves this goal. For non-supersymmetric models for Dirac leptogenesis, see ref. [5].

Our spectrum extends from that of the Minimal Supersymmetric Standard Model [6] only in the sector that is a singlet under the Standard Model gauge group. This extension is in the spirit of the nearly Minimal Supersymmetric Standard Model nMSSM [7]. In the latter, one adds to the MSSM superpotential just one term, coupling a singlet chiral superfield  $S$  to the two Higgs superfields of the MSSM, in order to generate the  $\mu$ -term. In so far as direct effects on the low energy superpotential are concerned, we need to add to the nMSSM spectrum the right chiral neutrino superfields  $N_i$ ,  $i$  standing for flavor, heavy Dirac pairs of singlet chiral superfields  $\Phi_k, \Phi_k^c$  carrying  $B - L$  charges  $\pm 1$ , and a generic heavy singlet superfield  $X$  mediating supersymmetry breaking. In order to quicken the decays of relic right-handed sneutrinos, we need to have large values of certain supersymmetry breaking parameters whose contributions to the scalar potential get scaled by the neutrino mass and remain small without much affecting the lightest sparticle spectrum. Such large supersymmetry breaking parameters may arise consistently from a dynamical supersymmetry breakdown in the hidden sector. We show this explicitly by constructing a specific scheme in which the effects of such a breaking are fed down to the observable (neutrino) sector by the field  $X$  associated with  $U(1)_m$  mediation [8, 9].

In section II we describe our mechanism for Dirac neutrino leptogenesis with a toy superpotential. In section III we discuss the required strengths of the supersymmetry breaking parameters to avoid the late non-thermal overproduction of LSPs. The origin of such parameters from a dynamically broken supersymmetry in the hidden sector, fed down by  $U(1)$  gauge mediation, is outlined for a particular scheme in section IV. Finally, section V summarizes our conclusions.

## 2. Mechanism for Dirac neutrino leptogenesis

Consider first the toy superpotential

$$W_{\text{Dirac}} = h_{ik} L_i H_2 \Phi_k^c + h'_{jk} N_j S \Phi_k + M_k \Phi_k \Phi_k^c, \tag{2.1}$$

which will later constitute part of our full superpotential for the observable sector. Here  $h, h'$  are Yukawa coupling strengths and  $M_k$  are heavy masses. We have used the standard notation for the superfields  $L_i, N_i, H_2$ , etc. Integrating out the heavy pairs, we have

$$W_{\text{eff}} = \frac{h_{ik} h'_{jk}}{M_k} L_i H_2 N_j S \equiv \frac{m_{ij}^\nu}{v_2 v_S} L_i H_2 N_j S, \tag{2.2}$$

generating the Dirac neutrino mass matrix  $m_{ij}^\nu$  with the vacuum expectation values (VEVs)  $v_2 = \langle H_2^0 \rangle$  and  $v_S = \langle S \rangle$ . Note that we need more than two pairs of heavy fields ( $\Phi_i, \Phi_i^c$ ) in order to generate at least two non-vanishing neutrino mass eigenvalues. In this approach, a tiny Dirac neutrino coupling  $y = m^\nu / v_2 \sim 10^{-13}$  is obtained from reasonable values of

the Yukawa coupling strengths  $hh' \sim 10^{-6}$  and the heavy mass scale  $M_i \sim 10^{10}$  GeV with  $v_S \sim \text{TeV}$ . At this point we do not specify the origin of the VEV and the mass associated with the singlet superfield  $S$ .

Let us assume two pairs of heavy superfields for simplicity, which gives us two non-vanishing neutrino mass eigenvalues. When the masses of the heavy neutrinos are hierarchical,  $M_1 \ll M_2$ , the decays of the lightest superfields  $\Phi_1, \bar{\Phi}_1^c$  will produce the asymmetries  $\epsilon$  and  $-\epsilon$  in the final states  $LH_2$  and  $\bar{N}\bar{S}$ , respectively. The decay rate of  $\Phi_1, \bar{\Phi}_1^c$  is

$$\Gamma_D = \frac{1}{8\pi} \sum_i [2|h_{i1}|^2 + |h'_{i1}|^2] M_1 \quad (2.3)$$

and the asymmetry  $\epsilon$ , defined by  $\epsilon \Gamma_D \equiv \Gamma(\Phi_1, \bar{\Phi}_1^c \rightarrow LH_2) - \Gamma(\bar{\Phi}_1, \Phi_1^c \rightarrow \bar{L}\bar{H}_2)$ , is

$$\epsilon = \frac{\sum_{i,j} \text{Im}(h_{i1} h_{i2}^* h'_{j1} h'_{j2}^*)}{4\pi \sum_k (2|h_{k1}|^2 + |h'_{k1}|^2)} \frac{\delta}{1 - \delta^2}, \quad (2.4)$$

where  $\delta \equiv M_1/M_2$  [4]. Here we assume that the heavy masses  $M_i$  are real and positive without loss of generality. For a successful leptogenesis, the lepton asymmetry, normalized by the entropy density  $s$ ,  $Y_L \equiv (n_L - n_{\bar{L}})/s$ , is required to be

$$Y_L \approx \epsilon Y_\Phi \sim 10^{-10}, \quad (2.5)$$

where  $Y_\Phi$  is the out-of-equilibrium density of the heavy superfield pair  $\Phi_1, \bar{\Phi}_1^c$ . This is usually expressed in terms of the ratio between the relativistic equilibrium number density and the entropy density  $n_\Phi^{\text{eq}}/s \sim 10^{-3}$  as well as the efficiency factor  $\eta \lesssim 1$ :  $Y_\Phi \equiv (n_\Phi^{\text{eq}}/s)\eta$ . The efficiency factor, which is controlled by the predominance of the decay over the inverse decay process, is determined by the ratio between the decay rate and the expansion parameter of the universe  $H(T)$  at  $T = M_1$ :  $K \equiv \frac{\Gamma_D}{H(M_1)} \approx \frac{0.1}{g_*^{1/2}} \frac{M_P}{M_1} \sum_k (2|h_{k1}|^2 + |h'_{k1}|^2)$  where  $M_P$  is the Planck mass and  $g_*$  is the relativistic degree of freedom at  $T = M_1$ . The approximate functional form of the efficiency factor can be expressed as  $\eta \approx 1/K(\ln K)^{0.8}$  in the case of  $K \gg 1$  [10]. From the relation (2.2), one typically gets  $K \sim (m_\nu/10^{-3}\text{eV})(\text{TeV}/v_S)$  leading to  $K \sim 10$  (and thus  $\eta \sim 0.1$ ) for the solar neutrino mass scale  $m_\nu \sim 0.01$  eV and  $v_S \sim \text{TeV}$ . A more specific choice of parameters for this will be shown below. To achieve the maximal efficiency  $\eta \sim 1$  corresponding to  $K \sim 1$ , we may take  $v_S \sim 10$  TeV or introduce one more pair of heavy fields associated with a smaller neutrino mass,  $m_\nu \sim 10^{-3}$  eV, as in the case of the usual seesaw mechanism [10].

Let us remark that the asymmetry (2.4) is suppressed by the factor  $\delta \ll 1$  for hierarchical heavy neutrino masses. Nevertheless, a realistic value of  $\epsilon$  can be attained [4, 10] by choosing somewhat large magnitudes for  $h, h'$ . On the other hand, the CP asymmetry can be resonantly enhanced if  $\delta \approx 1$ . This possibility with  $h \sim h'$  may well be motivated to generate the observed mild hierarchy in the light neutrino masses. In this case, the CP asymmetry is generalized to

$$\epsilon = \sum_l \frac{\sum_{i,j} \text{Im}(h_{i1} h_{i2}^* h'_{j1} h'_{j2}^*)}{4\pi \sum_k (2|h_{kl}|^2 + |h'_{kl}|^2)} \frac{\delta}{1 - \delta^2}, \quad (2.6)$$

taking into account the contributions from the decays of two pairs  $\Phi_l, \Phi_l^c$  with  $l = 1, 2$ . The above formula can be straightforwardly generalized to the case of three families of heavy superfields. For a detailed analysis of the parameter region satisfying the leptogenesis conditions, we refer the readers to refs. [4, 10]. As an example of the parameters accommodating a successful leptogenesis, let us choose  $M_i \sim 10^{10}$  GeV and  $h \sim h' \sim 10^{-3}$ . This gives us  $K \sim 10$  leading to  $\eta \sim 0.1$  and thus  $Y_\Phi \sim 10^{-4}$  [10]. Now taking  $1 - \delta \sim 0.1$ , we get the desired value of  $\epsilon \sim 10^{-6}$ . Note also that one can have almost degenerate heavy fields,  $\delta \simeq 1$ , in which case the condition (2.5) can be met by significantly enhanced  $\epsilon$  and much smaller  $Y_\Phi$ .

Independently of the details of the parameter space, one can conclude that successful leptogenesis via either (2.4) or (2.6) implies  $Y_\Phi > 10^{-10}$ . This will lead in a supersymmetric scenario to the problem of unwanted relics which we propose to solve below.

### 3. Conditions on supersymmetry breaking parameters

A key requirement of Dirac leptogenesis is not to equilibrate the right-handed neutrinos with left-handed ones both above and at the electroweak phase transition when sphaleron interactions are active. This is needed to make the asymmetry in the left-handed neutrino turn into the baryon asymmetry via those interactions [2]. In other words, the scattering rates  $\Gamma_S$ , induced by the effective operator (2.2), needs to be suppressed as compared with the expansion rate  $H$  of the universe above and at the electroweak phase transition:  $\Gamma_S < H$  for  $T \geq T_c$ , where  $T_c \sim 100$  GeV. For the Dirac neutrino Yukawa operator  $W = (m_\nu/v_2) LH_2N$ , this condition requires  $m_\nu \lesssim 10$  keV [2] which is trivially satisfied with the observed neutrino mass scale  $0.05 \text{ eV} \lesssim m_\nu \lesssim 0.33 \text{ eV}$ . Similarly, one can find that the non-equilibration of the effective operator (2.2) itself is satisfied for a temperature below  $T_\nu$ :

$$T_\nu \sim 4 \times 10^{15} \text{ GeV} \left( \frac{0.05 \text{ eV}}{m_\nu} \right)^2 \left( \frac{v_S}{1 \text{ TeV}} \right)^2, \tag{3.1}$$

which puts the bound;  $M \lesssim T_\nu$ , or  $v_S \gtrsim 1.6$  GeV for  $m_\nu \gtrsim 0.05$  eV and  $M \sim 10^{10}$  GeV.

The same argument has to be applied to the supersymmetry breaking operators associated with the supersymmetric operators (2.2). Taking general supersymmetry breaking parameters in the scalar potential  $V$ , we write

$$V = \frac{m_\nu F_S}{v_2 v_S} \tilde{l} h_2 \tilde{n} + \frac{\Lambda_\nu m_\nu}{v_2 v_S} \tilde{l} h_2 \tilde{n} s, \tag{3.2}$$

where  $F_S$  is the  $F$ -term of  $S$  and  $\Lambda_\nu$  is a supersymmetry breaking parameter which can come from the usual  $A$ -term or  $B$ -term of the heavy field mass operator, as will be shown below. For the moment, we do not specify the origins of these supersymmetry breaking parameters keeping their values arbitrary. The requirement of not equilibrating the above operators (3.2) leads to the conditions

$$\frac{F_S}{v_S} \lesssim 5 \times 10^7 \text{ GeV} \quad \text{and} \quad \frac{\Lambda_\nu}{v_S} \lesssim 5 \times 10^5. \tag{3.3}$$

With these bounds, we find the couplings in eq. (3.2) satisfying  $m_\nu F_S/v_2 v_S \lesssim 10^{-5}$  GeV and  $m_\nu \Lambda_\nu/v_2 v_S \lesssim 10^{-7}$ . These values are too small (compared with the weak scale) to lead to any observable consequences on sparticle spectra and couplings probed in collider experiments. In particular, it is impossible to realize the mixed sneutrino dark matter [11, 12] in the framework of Dirac leptogenesis.

Another important issue in supersymmetric Dirac leptogenesis is the problem with unwanted relics. During leptogenesis, the out-of-equilibrium decays of the heavy fields  $(\Phi_i, \Phi_i^c)$  produce the scalar components of  $N$  and  $S$  by the amount of  $Y_X$ . Since their couplings are small as in eqs. (3.2, 3.3), they never get equilibrated. Thus their initial abundances are retained until they decay to the lightest supersymmetric particle (LSP) which is considered to be dark matter. If such a non-thermally produced LSP survives, the observed dark matter abundance today would put the upper bound

$$Y_X m_X \lesssim 10^{-10} \text{ GeV}. \tag{3.4}$$

This implies  $Y_X \lesssim 10^{-12}$  for the LSP mass around 100 GeV which is clearly in contradiction to the condition (2.7) for successful leptogenesis. To avoid this difficulty, we need to have large couplings in eq. (3.2) to make the decay of the scalar  $\tilde{n}$  occur while the LSP can equilibrate. Requiring the decay temperature  $T_D$  to be larger than the LSP decoupling temperature  $T_{\text{LSP}}$ , which we take to be 10 GeV, we find

$$\Lambda_\nu \text{ or } \frac{F_S}{v_S} \gtrsim 3 \times 10^6 \left( \frac{\tilde{m}_N}{\text{TeV}} \right)^{\frac{3}{4}} \left( \frac{T_{\text{LSP}}}{10 \text{ GeV}} \right)^{\frac{3}{4}} \text{ GeV}. \tag{3.5}$$

Similar bounds emerge from the scalar  $s$  decay. This consideration is also applicable when the scalar  $\tilde{n}$  (or  $s$ ) is the LSP. From eqs. (3.3, 3.5), the allowed range of the supersymmetry breaking parameters is found to be  $\Lambda_\nu$  or  $F_S/v_S \approx (10^6 - 10^8)$  GeV. The challenge then is to generate such a large supersymmetry breaking parameter. It should be noted again that the corresponding contribution to the scalar potential  $V$  in eq. (3.2), being scaled by the tiny neutrino mass, is controllably small and is not expected to give rise to any undesirably large sparticle mass or any other significant phenomenological consequence at laboratory energies. In the next section, we will present a consistent framework for Dirac leptogenesis realizing the conditions  $F_S/v_S < \text{TeV}$  but  $\Lambda_\nu \sim 10^{7-8}$  GeV with  $v_S \sim \text{TeV}$ .

#### 4. Origin of $v_S$ and $\Lambda_\nu$

We now show that the conditions for successful Dirac leptogenesis described above can be realized consistently in the context of a suitably extended version of nMSSM. As in the usual version, let us first couple the singlet superfield  $S$  to the Higgs mass operator through the superpotential term  $SH_1 H_2$  and thus generate  $v_S \sim \text{TeV}$  dynamically. In this scheme, there is no problem with non-thermally produced  $S$  as it can equilibrate or decay fast by its large coupling with Higgs fields. For a viable scenario of nMSSM containing the neutrino sector (2.1), we introduce the superfields and two global symmetries R and B-L on top of MSSM as in table 1.

	$S$	$H_1$	$H_2$	$L$	$N$	$\Phi$	$\Phi^c$	$X$
$R$	$-\frac{4}{3}$	$\frac{5}{3}$	$\frac{5}{3}$	$-\frac{1}{3}$	$\frac{8}{3}$	$\frac{2}{3}$	$\frac{2}{3}$	$\frac{2}{3}$
$B - L$	0	0	0	1	-1	-1	1	0
$Y$	0	$-\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	0	0	0	0

**Table 1:** The symmetry assignment of our model.

Here we have taken the R-charge normalization such that  $R(W) = 2$  and  $Y$  as the Standard Model hypercharge. This allows us to write down the superpotential

$$W = \lambda S H_1 H_2 + h_{ik} L H_2 \Phi_k + h'_{ik} N S \Phi_k^c + k_1 X \Phi \Phi^c + \frac{k_2}{3} X^3 + \frac{k_3}{5} \frac{X^5}{M_P^3} S. \quad (4.1)$$

It is now crucial for us to generate high scales for  $v_X = \langle X \rangle$  and  $F_X$ , which can be easily realized by adopting the idea of dynamical supersymmetry breaking fed down by gauge messengers [8, 9]. As an illustration, let us consider the  $SU(6) \times U(1) \times U(1)_m$  model proposed in ref. [9]. This model contains chiral superfields with the quantum numbers as follows:

$$\mathcal{A}(15, 1, 0), \bar{\mathcal{F}}^\pm(\bar{6}, -2, \pm 1), \mathcal{S}^0(1, 3, 0), \mathcal{S}^{\pm 2}(1, 3, \pm 2).$$

Without contradicting our symmetry assignments in table 1, we can extend the superpotential (4.1) to include the terms in eqs. (2.19), (2.23) and (3.1) of ref. [9]. Thus our dynamical supersymmetry breaking (DSB) is described by the superpotential

$$W_{\text{DSB}} = \frac{\lambda_6}{M_P} \mathcal{A} \bar{\mathcal{F}}^+ \bar{\mathcal{F}}^- \mathcal{S}^0 + \frac{\Lambda_6^5}{(\mathcal{A} \bar{\mathcal{F}}^+ \bar{\mathcal{F}}^- \mathcal{A}^3)^{1/3}} + k_4 X \phi^+ \phi^- + \frac{k_2}{3} X^3, \quad (4.2)$$

where  $\lambda_6$  is a Yukawa coupling,  $\Lambda_6$  is the dynamical scale of  $SU(6)$ ,  $\phi^\pm$  are the messenger fields carrying the messenger hypercharge  $\pm 1$  under  $U(1)_m$ , and  $X$  is the heavy singlet superfield used in eq. (4.1). The strong dynamics of  $SU(6)$  generates the second (non-perturbative) term which drives dynamical supersymmetry breaking in combination of the first (tree-level) term. Unlike the original propose for the gauge mediated supersymmetry breaking [9], we will assume a high scale  $\Lambda_6 \sim 10^{13}$  GeV so that the dynamically generated  $F$ -term  $F_G \sim \lambda_6^{1/2} \Lambda_6^{5/2} / M_P^{1/2}$  induces  $\tilde{m} \sim F_G / M_P \sim \text{TeV}$ , which is a right value for the gravitino mass and all the soft masses of the MSSM through gravity mediation. On the other hand, dynamical supersymmetry breaking generates a negative mass-squared for the  $U(1)_m$ -charged scalar field  $\phi^\pm$ , and thereby  $v_X$  and  $F_X$  are induced through the last two terms of eq. (4.2) as in ref. [8, 9].

Now we can show explicitly how the conditions (3.3), (3.5) are met in our scheme. From the superpotential terms (4.1), one finds

$$M = k_1 v_X, \quad \Lambda_\nu = \frac{F_X}{v_X} \quad \text{and} \quad F_S = \frac{k_3}{5} \frac{v_X^5}{M_P^3}. \quad (4.3)$$

One can easily adjust the Yukawa couplings in eqs. (4.1), (4.2) to get, for instance,

$$v_X \sim 10^{11} \text{ GeV} \quad \text{and} \quad F_X / v_X \sim 10^8 \text{ GeV}, \quad (4.4)$$



from which one obtains  $\Lambda_\nu \sim 10^8$  GeV and  $F_S \ll \tilde{m}^2$ . The latter is too small to become relevant in generating the vacuum expectation values of the Higgs bosons and scalar  $S$ . However, we can have in the scalar potential

$$V = k_3 \frac{v_X^4 F_X}{M_P^3} s \sim \tilde{m}^3 s \tag{4.5}$$

which play the role of a tadpole contribution in the nMSSM [7]. This term, together with the coupling  $SH_1H_2$ , gives rise to  $v_S \sim \text{TeV}$  and thus appropriate  $\mu$  and  $B\mu$  terms respecting the Higgs stability and electroweak symmetry breaking conditions [6]. Let us finally note that  $\Lambda_\nu$  is in fact the  $B$ -term of the heavy mass operator  $M\Phi\Phi^c$ . This makes it clear why we introduced heavy singlets instead of  $SU(2)$  doublets. If the heavy fields ( $\Phi, \Phi^c$ ) were  $SU(2)$  doublets as in ref. [3],  $\Lambda_\nu$  would induce via gauge mediation [8, 9] too large soft supersymmetry breaking terms in the MSSM sector characterized by  $\tilde{m} \sim \alpha_2 \Lambda_\nu / 4\pi$ .

## 5. Conclusion

We have successfully implemented Dirac leptogenesis in an extended version of the nMSSM. Unlike that of Murayama and Pierce [3], our extension of the MSSM involves only superfields which are singlets under the Standard Model gauge group. The advantage here is the quickening of the otherwise undesirably slow decays of relic singlet sneutrinos, thereby avoiding a non-thermal overproduction of LSPs which had plagued earlier Dirac leptogenesis scenarios. This is achieved by generating heavy mass scales  $M \sim 10^{10}$  GeV and large supersymmetry breaking parameters  $\Lambda_\nu$  in the range  $10^7 - 10^8$  GeV by dynamical supersymmetry breaking.

Apart from the usual signatures of the Higgs sector extended with singlet fields, our model has no signatures in low energy phenomenology. In particular, the contribution of the seemingly large parameter  $\Lambda_\nu$  to the scalar potential get scaled by the neutrino mass and is inconsequentially small in terms of probing the scalar right-handed neutrino sector. A clear laboratory distinction between our model and that for standard Majorana leptogenesis is that our light neutrinos are Dirac particles with lepton number conserving interactions. Thus, non-observation of neutrinoless nuclear double beta decaying the future experiments may hint at the Dirac nature of light neutrinos, providing indirect support for Dirac leptogenesis.

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