UCRHEP-T423
November 2006

Low-Energy Thermal Leptogenesis in an Extended NMSSM Model

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Abstract

Thermal leptogenesis in the canonical seesaw model in supersymmetry suffers from the incompatibility of a generic lower bound on the mass scale of the lightest right-handed neutrino and the upper bound on the reheating temperature of the Universe after inflation. This is resolved by adding an extra singlet superfield, with a discrete Z_2 symmetry, to the NMSSM (Next to Minimal Supersymmetric Standard Model). This generic mechanism is applicable to any supersymmetric model for lowering the scale of leptogenesis.

PACS numbers: 12.60.Jv; 14.60.Pq; 98.80.Cq

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I. INTRODUCTION

In the minimal standard model (SM), neutrinos are massless. However, small nonzero neutrino masses are required by the atmospheric and solar neutrino experiments. A natural explanation for such tiny neutrino masses in the SM comes from an effective dimension-5 operator [1]

$$\mathcal{L}_\Lambda = \frac{f_{\alpha\beta}}{\Lambda}(\nu_\alpha\phi^0 - l_\alpha\phi^+)(\nu_\beta\phi^0 - l_\beta\phi^+) + H.c., \tag{1}$$

where $(\nu_\alpha, l_\alpha),\ \alpha = e, \mu, \tau$ are the usual left-handed lepton doublets transforming as $(2, -1/2)$ under the standard electroweak $SU(2)_L \times U(1)_Y$ gauge group and $(\phi^+, \phi^0) \sim (2, 1/2)$ is the usual Higgs doublet of the SM. There are three realizations of this operator [2], the most popular one being the canonical seesaw [3] mechanism which adds three singlet heavy neutral fermions $N_i,\ i = 1, 2, 3$ to the SM Lagrangian. The neutrino mass matrix is then given by

$$\mathcal{M}_\nu = -\mathcal{M}_D\mathcal{M}_N^{-1}\mathcal{M}_D^T \tag{2}$$

where $\mathcal{M}_D$ is the 3×3 Dirac mass matrix linking ν_α with N_i through the Yukawa interactions $h_{\alpha i}(\nu_\alpha\phi^0 - \ell_\alpha\phi^+)N_i$.

The Majorana masses of N_i violate lepton number by two units. Therefore, in the early Universe, a net lepton asymmetry may be generated [4] through the out-of-equilibrium decay of the lightest N_i (call it N_1). The generated lepton asymmetry then gets converted into a baryon asymmetry through the interactions of the SM sphalerons [5] which conserve $B - L$, but violate $B + L$, where B and L are baryon and lepton number respectively. The existence of N_i explains thus at the same time why both neutrino masses as well as the observed baryon asymmetry of the Universe (BAU) are nonzero and small.

In supersymmetric theories the reheating temperature (T_h) following inflation is likely to be rather low [6, 7]. Although some models [8] may allow a higher reheating temperature, in the conventional models T_h is bounded strongly from above by the possible overproduction of gravitinos. On the other hand, in the simplest version of the seesaw mechanism the condition for thermal leptogenesis requires that the mass M_1 of the lightest right-handed neutrino N_1 should be much higher than T_h, assuming that N_1 contributes to the left-handed neutrino masses dominantly. Since inflation would erase any pre-existing lepton asymmetry, the asymmetry generated by N_1 after inflation would be highly suppressed by its small number density and hence this mechanism will fail to explain the BAU.

To avoid this problem, several ideas have been discussed in the literature [10, 11, 12]. An attractive scenario is the extended seesaw mechanism [13, 14]. In this paper we follow the same scheme and work with the canonical seesaw mechanism (SM plus three N_i) in the Next to Minimal Supersymmetric Standard Model (NMSSM), which has an extra singlet superfield χ. To distinguish N_i from χ, an exactly conserved Z_2 discrete symmetry is imposed, corresponding to $(-1)^L$. We then add an extra heavy singlet superfield S, together with a softly broken discrete symmetry Z_2', under which S is odd and all others are even. As a result, the production and decay channels of S are different, and the out-of-equilibrium decay of S can take place much below the mass scale of the lightest right-handed neutrino N_1. The lower bound on the mass scale of S can then be compatible with the upper bound on the reheating temperature after inflation.

The rest of this paper is arranged as follows. In section II we review the canonical leptogenesis and briefly recall the Davidson-Ibarra (DI) bound on the mass scale of N_1. In section III we discuss an extended seesaw model by introducing an additional heavy singlet fermion S of mass less than that of N_1. In section IV we discuss how the thermal-leptogenesis constraint on the mass scale of S can be lowered in comparison with the mass scale of N_1. In section V we solve the required Boltzmann equations numerically and show how the low mass scale of S is compatible with thermal leptogenesis. In section VI we state our conclusions.

II. CANONICAL LEPTOGENESIS AND DI BOUND

In canonical leptogenesis the lightest right-handed neutrino N_1 decays into either $\ell^-\phi^+$ and $\nu\phi^0$, or $\ell^+\phi^-$ and $\bar{\nu}\bar{\phi}^0$. Thus a CP asymmetry can be established from the interference of the tree-level amplitudes with the one-loop vertex [4] and self-energy corrections [15]. A net lepton asymmetry arises when the decay rate

$$\Gamma_1 = \frac{1}{8\pi v^2} \tilde{m}_1 M_1^2 \tag{3}$$

fails to compete with the expansion rate of the Universe

$$H(T) = 1.66 g_*^{1/2} \frac{T^2}{M_{pl}} \tag{4}$$

at $T \sim M_1$, where $\tilde{m}_1 = (m_D^\dagger m_D)_{11}/M_1$ is the effective neutrino mass parameter, $g_* \simeq 228$ is the effective number of relativistic degrees of freedom in the MSSM and $M_{pl} = 1.2 \times 10^{19}$ GeV. This means that a <u>upper</u> bound on $\tilde{m}_1$ may be established by first considering the out-of-equilibrium condition $H(T = M_1) > \Gamma_1$ which gives

$$\tilde{m}_1 < 1.6 \times 10^{-3} eV\ . \tag{5}$$

However, for $\tilde{m}_1 > 10^{-3} eV$ a reduced lepton asymmetry may still be generated, depending on the details of the Boltzmann equations which quantify the deviation from equilibrium of the process in question.

Assuming a normal mass hierarchy in the right handed neutrino mass spectrum the CP asymmetry ϵ_1 is given by

$$\epsilon_1 \simeq -\frac{3}{8\pi v^2} \left(\frac{M_1}{M_2}\right) \frac{Im[(m_D^\dagger m_D)_{12}]^2}{(m_D^\dagger m_D)_{11}}\ . \tag{6}$$

The baryon-to-photon ratio of number densities has been measured [16] with precision, i.e.

$$\eta_B \equiv \frac{n_B}{n_\gamma} = 6.1^{+0.3}_{-0.2} \times 10^{-10}. \tag{7}$$

To get the correct value of η_B, one needs ϵ_1 to be of order 10^{-6} to 10^{-7}. However, using the DI upper bound on the CP asymmetry parameter [9]

$$|\epsilon_1| \leq \frac{3M_1}{8\pi v^2} \sqrt{\Delta m^2_{atm}} \tag{8}$$

one can get a lower bound on the mass scale of the lightest right-handed neutrino as

$$M_1 \geq 2.9 \times 10^9 GeV \left(\frac{\eta_B}{6.1 \times 10^{-10}}\right) \left(\frac{4 \times 10^{-3}}{(n_{N_1}/s)\delta}\right) \left(\frac{v}{174 GeV}\right)^2 \left(\frac{0.05 eV}{\sqrt{\Delta m^2_{atm}}}\right)\ , \tag{9}$$

where v is the vacuum expectation value (vev) of the SM Higgs and it is also assumed that the physical left handed neutrinos follow the normal mass hierarchy. Since T_h is not likely to exceed 10^9 GeV, this poses a problem for canonical leptogenesis. In order to overcome this problem we consider an extended seesaw model as follows.

III. THE MODEL FOR THERMAL LEPTOGENESIS BELOW THE DI BOUND

In a recent paper [13] we proposed a singlet mechanism to overcome the DI bound shown in Eq. (9). We now present a realistic model where this mechanism can be implemented. One more ingredient has been added to produce these singlet fields abundantly in a thermal bath.

We start with the NMSSM model, which includes a singlet superfield χ in addition to the usual particles of the Minimal Supersymmetric Standard Model (MSSM). To implement the seesaw mechanism, we also include three right-handed neutrinos $N_i,\ i = 1, 2, 3$. We then demonstrate how a minimal extension of this model may admit a very low scale of leptogenesis, thus overcoming the gravitino problem. We include another singlet superfield S and impose a $Z_2 \times Z_2'$ discrete symmetry. Under Z_2, the lepton superfields L_i, l_i^c, N_i, S are odd, whereas the Higgs superfields $\phi_{1,2}, \chi$ are even. This corresponds to having an exactly conserved lepton number $(-1)^L$, or the usual $R-$parity of the MSSM. Under Z_2', S is odd and all others are even, but Z_2' is allowed to be broken softly. The most general superpotential invariant under $Z_2 \times Z_2'$ is then given by

$$\begin{aligned} W \ &= \ h^e_{ij} L_i l^c_j \phi_1 + h_{ij} L_i N_j \phi_2 + \mu\phi_1\phi_2 + M_{ij} N_i N_j + M_\chi \chi\chi \\ &\quad + \alpha\chi\chi\chi + \beta\chi\phi_1\phi_2 + f_N \chi N_i N_j + M_S SS + f_S \chi SS. \end{aligned} \tag{10}$$

We do not discuss quarks nor other interactions of the NMSSM, which have been studied elsewhere. We deal only with neutrino masses and leptogenesis.

We now break Z_2' softly and the only possible such term is

$$W_s = d_i N_i S,$$

i.e. exactly as required to implement the singlet mechanism of ref. [13]. This allows S to mix with N_i to form a 7×7 mass matrix in the basis $[L_i \;\; N_i \;\; S]$, i.e.

$$\mathcal{M} = \begin{pmatrix} \mathbf{0} & m_D & 0 \\ m_D & M_N & d \\ 0 & d & M_S \end{pmatrix} \tag{11}$$

where $M_S = f_S \langle \chi \rangle$, and without loss of generality we choose M_N to be diagonal with masses $M_{1,2,3}$. For small $d_i/(M_i - M_S)$ as well as the usual seesaw assumption that the entries of m_D are very small relative to M_i, the heavy masses are roughly given by

$$\begin{aligned} M_{S'} &\simeq M_S - \sum_i \frac{d_i^2}{M_i - M_S}, \\ M_{N_i'} &\simeq M_i + \frac{d^2}{M_i - M_S}, \end{aligned} \tag{12}$$

corresponding to the mass eigenstates S' and N_i'

$$\begin{aligned} S' &\simeq S - \sum_i \frac{d_i}{M_i - M_S} N_i \\ N_i' &\simeq N_i + \frac{d_i}{M_i - M_S} S. \end{aligned} \tag{13}$$

The light neutrino mass matrix is then

$$(m_\nu)_{ij} \simeq - \sum_k (m_D)_{ik} \left(M_k + \frac{d_k^2}{M_k - M_S} \right)^{-1} (m_D)_{kj}. \tag{14}$$

In the limit $d_i \to 0$ we recover the neutrino masses as in the canonical seesaw mechanism. We assume thus $M_{S'} \simeq M_S$ and $M_{N_i'} \simeq M_N$ in the following.

IV. LEPTON ASYMMETRY AND LOWER BOUND ON M_S

In this model, the addition of S allows the choice $M_S < M_1$. The induced couplings of S to leptons are suppressed by factors of d_i/M_i compared to those of N_i. The decay rate of S is thus given by

$$\Gamma_D^S = \frac{1}{8\pi v^2} M_S \sum_i \left[\tilde{m}_i M_i (d_i/M_i)^2 \right] , \tag{15}$$

where we have assumed $M_s \ll M_i$ and the effective neutrino mass parameter is defined as

$$\tilde{m}_i = \frac{(m_D^\dagger m_D)_{ii}}{M_i} \,. \tag{16}$$

Assuming that $\frac{d_1}{M_1} = \frac{d_2}{M_2} = \frac{d_3}{M_3}$ and $\tilde{m}_3 > \tilde{m}_2 > \tilde{m}_1$ the above Eq. (15) can be rewritten as

$$\Gamma_D^S \simeq \frac{1}{8\pi v^2} M_S \tilde{m}_3 M_3 (d_3/M_3)^2 \,. \tag{17}$$

The out-of-equilibrium condition $\Gamma_D^S < H(T \sim M_S)$ is thus suppressed by a factor

$$\eta = \left(\frac{d_3^2}{M_3 M_S}\right)\left(\frac{\tilde{m}_3}{\tilde{m}_1}\right) \equiv \kappa \left(\frac{\tilde{m}_3}{\tilde{m}_1}\right) \tag{18}$$

in comparison to $\Gamma_1/H(T \sim M_1)$ and hence can be satisfied at a lower mass depending on the value of d_3. The value of $\tilde{m}_3$, $\tilde{m}_1$ and M_3 can be approximately fixed from the low energy neutrino oscillation data. So, the remaining free parameters are d_3 and M_S on which the suppression factor η depends.

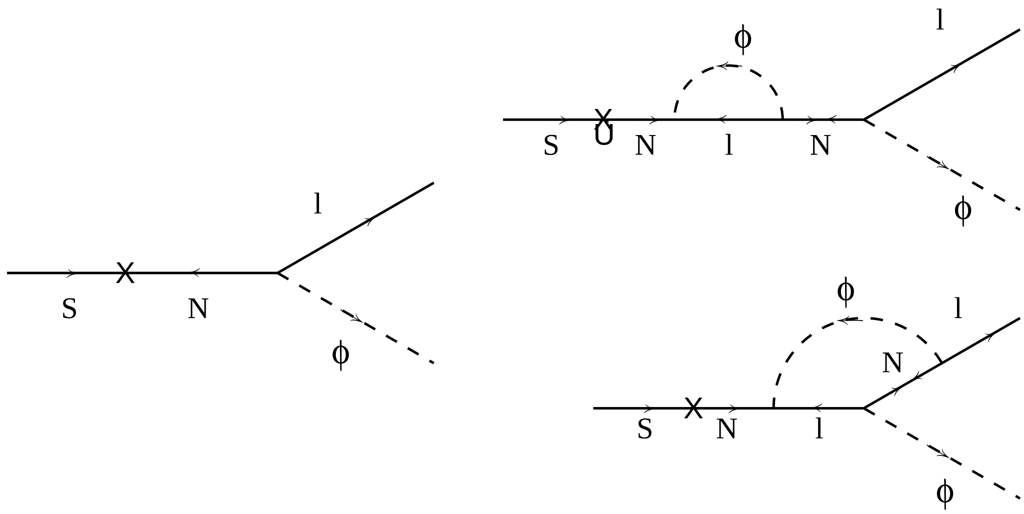

FIG. 1: Tree-level and one-loop (self-energy and vertex) diagrams for S decay, which interfere to generate a lepton asymmetry.

The CP asymmetry generated by the decays of S comes from the interference of the tree-level and one-loop diagrams of fig. (1). Both the numerator and denominator of Eq. (6) are then suppressed by the same $(d_i/M_i)^2$ factor, and we obtain

$$\epsilon_S \simeq -\frac{3}{8\pi v^2}\left(\frac{M_S}{M_2}\right)\frac{Im[(m_D^\dagger m_D)_{12}]^2}{(m_D^\dagger m_D)_{11}} \,, \tag{19}$$

where we have assumed that $M_1 \ll M_2 \ll M_3$. Thus we see that the CP asymmetry parameter is independent of the suppression parameter d_3. Therefore, depending on the value of κ the L-asymmetry will saturate at different temperatures as implied by Eq. (18). This is shown in Section V by numerically solving the required Boltzmann equations. As demonstrated in fig. (8) the value of M_S can be lowered much less than the DI bound on M_1 by an appropriate choice of d_3. This is because the low values of M_S are not restricted by the low energy neutrino oscillation data for $d_i \ll M_i$ as we have seen from Eq. (12). Moreover the washout effects are suppressed for low values of (d_3/M_3). So, a successful lepton asymmetric universe before the electroweak phase transition can be created even for a TeV scale of S.

V. NUMERICAL ESTIMATION OF LEPTON ASYMMETRY

A. Production and decay of S

In this model S is produced through the decay of χ. The corresponding Yukawa coupling f_S can be as large as of order unity. Therefore, S can be brought to thermal equilibrium through the scattering processes: $S\bar{S} \to S\bar{S}$, $S\bar{S} \to \chi\chi^\dagger$ and $\chi S \to \chi S$. Note that these processes never change the number density of S, but they keep S in kinetic equilibrium. The decay rate of χ can be given by

$$\Gamma_D^\chi = \frac{f^\dagger f}{8\pi} M_\chi \left(1 - \frac{4M_S^2}{M_\chi^2}\right)^{3/2} \frac{K_1\left((M_\chi/M_S)z\right)}{K_2\left((M_\chi/M_S)z\right)} \,. \tag{20}$$

where (K_1/K_2) is the boost factor and $z = M_S/T$. Thus the inverse decay of χ is given by $\Gamma_{ID}^\chi = (n_\chi^{eq}/n_S^{eq})\Gamma_D^\chi$.

Once the S particles are produced, they decay through the channel: $S \to \ell\phi^\dagger, \bar{\ell}\phi$ as shown in fig. (1) which violates lepton number by two units. Apart from that, the other process which depletes the number density of S is $S\ell \to \phi \to Q\bar{t}$. This is shown in fig. (2).

The subsequent decay of S, below its mass scale, then produces the required baryon asymmetry through the leptogenesis route. However, an exact lepton asymmetry can be

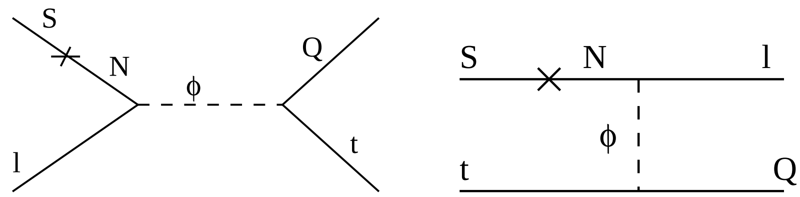

FIG. 2: $\Delta L = \pm 1$ processes which deplete the number density of S. These processes also deplete the net lepton number density produced through the decay channel.

estimated by solving the required Boltzmann equations [17]. It is useful to define the Boltzmann equations in terms of the dimensionless variables $Y_S = n_S/s$ and $Y_L = n_L/s$, where Y_S is called the comoving density of S while Y_L is the density of net lepton in a comoving volume and

$$s = \frac{2\pi^2}{45} g_* T^3 \tag{21}$$

is the entropy density. The required Boltzmann equations are given as

$$\frac{dY_S}{dz} = -\left(Y_S - Y_S^{eq}\right)\left[\frac{\Gamma_D^S}{zH(z)} + \frac{\Gamma_s^S}{zH(z)}\right] \tag{22}$$

and

$$\frac{dY_L}{dz} = \epsilon_S \frac{\Gamma_D^S}{zH(z)}\left(Y_S - Y_S^{eq}\right) - \frac{\Gamma_W^\ell}{zH(z)} Y_L\,, \tag{23}$$

where Γ_D^S, Γ_s^S and Γ_W^ℓ simultaneously represent the decay, scattering and wash out rates that takepart in establishing a net lepton asymmetry in a thermal plasma. The Hubble parameter $H(z)$ is given by

$$H(z) = \frac{H(M_S)}{z^2} \qquad \text{with} \qquad H(M_S) = 1.67 g_*^{1/2} \frac{M_S^2}{M_{pl}}\,. \tag{24}$$

In a relativistic frame the decay rate (17) can be rewritten as

$$\Gamma_D^S = \frac{1}{8\pi v^2} M_S \left(\frac{K_1(z)}{K_2(z)}\right) \tilde{m}_3 M_3 (d_3/M_3)^2\,. \tag{25}$$

The Γ_s^S in Eq. (22) represents the processes which violate lepton number by one unit and

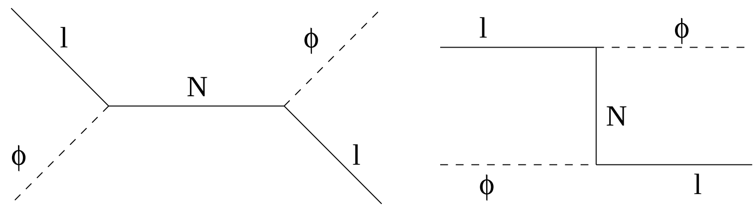

FIG. 3: $\Delta L = \pm 2$ processes which deplete the number density of net leptons.

is given by [1]

$$\Gamma_S = 4\Gamma^S_{\phi,s} + 2\Gamma^S_{\phi,t} \,. \tag{26}$$

The Γ_W in Eq. (23) represents the lepton number violating processes by two units and is given by

$$\Gamma_W = \frac{1}{2}\Gamma^S_{ID} + 2\Gamma^l_{\phi,t} + 2\Gamma^l_{\phi,s}\left(\frac{Y_{N1}}{Y^{eq}_{N1}}\right) + 2\Gamma^l_{N1} + 2\Gamma^l_{N1,t} \,, \tag{27}$$

where $\Gamma^S_{ID} = (n^{eq}_S/n^{eq}_l)\Gamma^S_D$. In Eqs. (26) and (27) the Γ's are defined as $\Gamma^x_i = (\gamma_i/n^{eq}_x)$ where γ_i is the scattering density. Note that the other $\Delta L = \pm 2$ processes: $ll \to N_1 S N_1 \to \bar{\phi}\bar{\phi}$, and of course higher order processes, which contribute to Γ^ℓ_W are suppressed in comparison to the processes shown in fig. (3).

B. Solution of Boltzmann equations

In fig. (4) we have plotted the decay and inverse decay of χ and S against z. It is shown that the inverse decay of $\ell + \phi^\dagger \to S$ is not sufficient to bring S into thermal equilibrium even if the suppression factor d_3 is as large as 10^8 GeV. On the contrary, the decay rate of χ is sufficiently larger than the Hubble expansion parameter. Hence S can be brought to thermal equilibrium through the scattering process involving χ as long as $(M_\chi/M_S) \simeq O(10^{1-2})$. Therefore, at a temperature far above the mass scale of S it is in thermal equilibrium and hence a net lepton asymmetry in the thermal plasma is zero. Below the mass scale of S

[1] We have not included the SUSY processes. It is shown that upon inclusion of those processes the result doesn't change significantly [18].

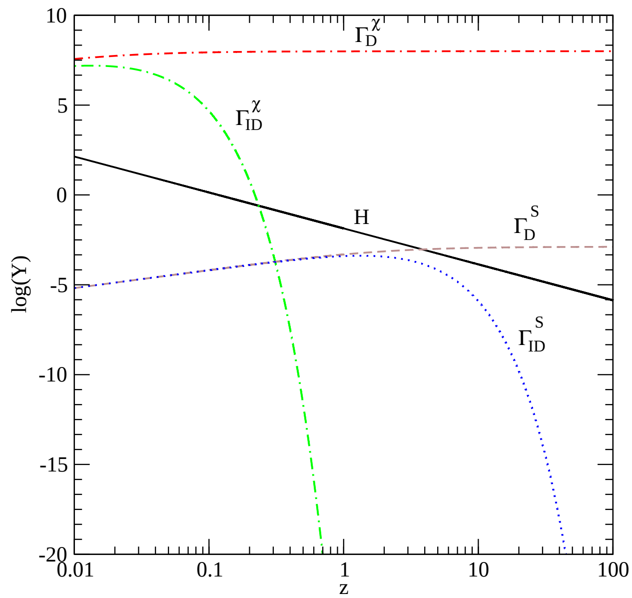

FIG. 4: The production and decay, and their inverse processes, are compared with the Hubble expansion parameter H for a typical set of parameters. We have used $M_S = 10^8$ GeV, $M_\chi = 10^{10}$ GeV, $M_3 = 10^{10} GeV$, $\tilde{m}_3 = 10^{-2}$ eV, $d_3 = 10^8$ GeV and $f_S = 0.5$.

the lepton number violating processes go out of thermal equilibrium and thus produce a net lepton asymmetry dynamically. This is obtained by solving the Boltzmann equations (22) and (23). We take the following initial conditions:

$$Y_S = Y_S^{eq} \quad \text{and} \quad Y_L = 0 \quad \text{at} \quad z \to 0 \,. \tag{28}$$

The evolution of the number density of S and the corresponding asymmetry with respect to z are shown in figs. (5), (6) and (7). At any epoch z the value of Y_S and the corresponding asymmetry Y_L can be inferred from

$$\frac{z}{sH(M_S)}\gamma_D^S \propto \kappa \tilde{m}_3 \,. \tag{29}$$

Since the decay rate of S is suppressed by a factor of κ, the asymmetry is produced at late times depending how small it is. However, the value of κ cannot be made indefinitely small. Because a net L-asymmetry has to build up before the electroweak phase transition which is required to be converted to the B-asymmetry through the sphaleron transitions. The

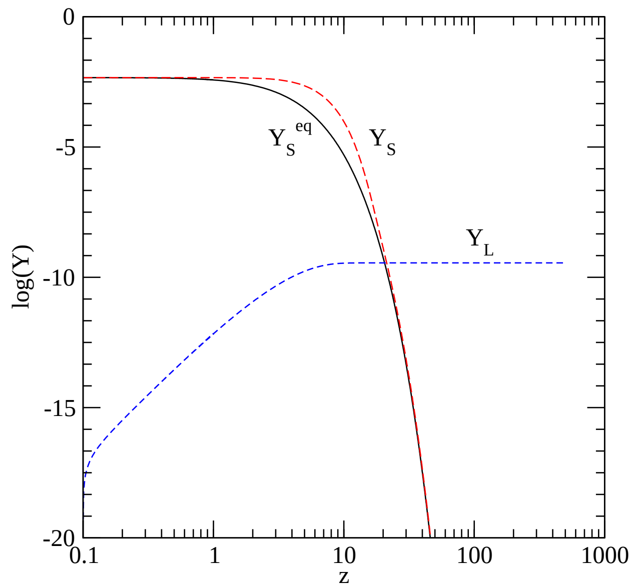

FIG. 5: The evolution of S is shown against z with $M_S = 10^8$ GeV, $M_1 = 10^9$ GeV, $M_3 = 10^{10}$ GeV and $d_3 = 10^8$ GeV and the CP asymmetry parameter is $\epsilon_S = 10^{-7}$.

final L-asymmetry is numerically obtained for three values of d_3 in figs. (5), (6) and (7). It is found that the final L-asymmetry is almost same apart from a numerical factor. This is because for the delayed decay of S the wash out effects are comparatively small. While $\kappa = 10^{-2}$ and 10^{-4} are used in figs. (5) and (6), it is of 10^{-6} in fig. (7). As seen from figs. (5), (6) and (7), for $\kappa = 10^{-2}, 10^{-4}, 10^{-6}$ the value of Y_L is saturated at around $z_s \simeq 10, 10^2, 10^3$ respectively. Assuming that a final L-asymmetry has to be produced before $T_{ew} \simeq 100$ GeV the minimum tolerable value of $\kappa = 10^{-12}$ is obtained. This indicates that for $M_3 = 10^{10}$ GeV, d_3^2 and M_S can be readjusted among themselves so as to get the suppression factor κ ranging from 10^{-2} to 10^{-12}. This is shown in fig.(8). The solid line in fig. (8) is obtained for $M_S = d_3$. The region above to that are defined by $d_3 > M_S$. So these values of d_3 are unnatural and are not allowed. While the region below to the solid line are defined by $d_3 < M_S$ and hence is allowed for naturalness. Thus we see that a wide range of M_S values from 10^3 GeV to 10^8 GeV are allowed that can produce the required lepton asymmetry before the electroweak phase transition.

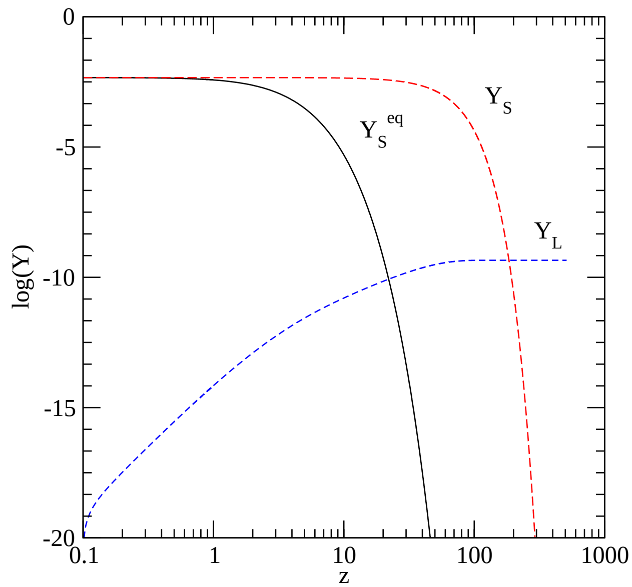

FIG. 6: The evolution of S is shown against z with $M_S = 10^8$ GeV, $M_1 = 10^9$ GeV, $M_3 = 10^{10}$ GeV and $d_3 = 10^7$ GeV and the CP asymmetry parameter is $\epsilon_S = 10^{-7}$.

VI. CONCLUSIONS

We accomplished the baryogenesis via leptogenesis from the decay of an additional singlet S in a supersymmetric extended NMSSM. The bound coming from the out-of-equilibrium condition could be evaded because the couplings of the singlets cancel out from the asymmetry, so the couplings could be small and can satisfy the out-of-equilibrium condition even at low scales. In the simplest seesaw models the couplings of the lightest right-handed neutrino could not be lowered much because that will not enable the thermal production of these fields. However, in the present case there is one additional singlet field (χ) which can produce these S fields, having large couplings to them, but itself not taking part in leptogenesis.

Acknowledgments

The work of EM was supported in part by the U. S. Department of Energy under Grant No. DE-FG03-94ER40837.

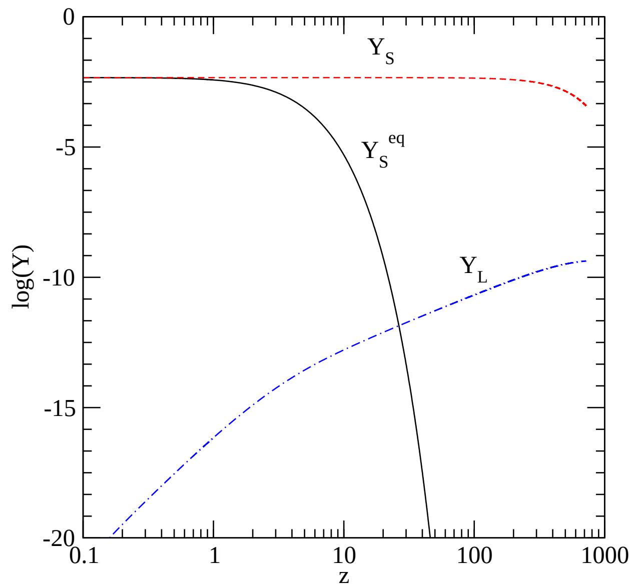

FIG. 7: The evolution of S is shown against z with $M_S = 10^8$ GeV, $M_1 = 10^9$ GeV, $M_3 = 10^{10}$ GeV and $d_3 = 10^6$ GeV and the CP asymmetry parameter is $\epsilon_S = 10^{-7}$.

APPENDIX A: SCATTERING DENSITIES

In this appendix we give the various scattering densities that have been used in the Boltzmann Eqs. (22) and (23).

$$\gamma_{\phi,s} = \frac{M_S^4 M_3 m_t^2}{64\pi^5 v^4 z}\left(\frac{d_3}{M_3}\right)^2 \tilde{m}_3 \int_1^\infty dx_1 \sqrt{x_1} K_1(z\sqrt{x_1}) \left[1 - \frac{1}{x_1}\right]^2 \tag{A1}$$

$$\gamma_{\phi,t} = \frac{M_S^4 M_3 m_t^2}{128\pi^5 v^4 z}\left(\frac{d_3}{M_3}\right)^2 \tilde{m}_3 \int_1^\infty dx_1 \sqrt{x_1} K_1(z\sqrt{x_1}) \left[1 - \frac{1}{x_1} + \frac{1}{x_1}\ln\left(\frac{x_1 - 1 + y}{y}\right)\right] \tag{A2}$$

where v is the vev of SM Higgs and $x_1 = \frac{s}{M_S^2}$, s being the Mandelstam variable, and $y = \frac{m_\phi^2}{M_s^2}$.

$$\gamma_{N1} = \frac{M_1^5 M_S \tilde{m}_1^2}{128\pi^5 v^4 z}\int_0^\infty dx_2 \sqrt{x_2} K_1\left(z\sqrt{x_2}\frac{M_1}{M_S}\right)$$
$$\left[1 + \frac{1}{D_1(x_2)} + \frac{x_2}{2D_1^2(x_2)}\left\{1 + \frac{1+x_2}{D_1(x_2)}\right\}\ln(1+x_2)\right] \tag{A3}$$

$$\gamma_{N1,t} = \frac{M_1^5 M_S \tilde{m}_1^2}{128\pi^5 v^4 z}\int_0^\infty dx_2 \sqrt{x_2} K_1\left(z\sqrt{x_2}\frac{M_1}{M_S}\right)\left[\frac{x_2}{x_2+1} + \frac{1}{x_2}\ln(x_2+1)\right] \tag{A4}$$

where $x_2 = \frac{s}{M_1^2}$ and

$$\frac{1}{D_1(x_2)} = \frac{x_2 - 1}{(x_2-1)^2 + \left(\frac{\Gamma_D^2}{M_1^2}\right)} \tag{A5}$$

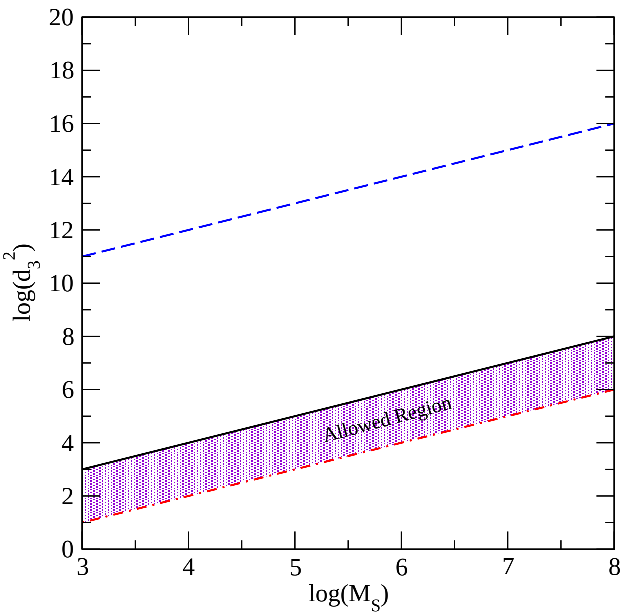

FIG. 8: The allowed values of M_S are shown in the plane of d_3^2 versus M_S.

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