Soft leptogenesis in the NMSSM with a singlet right-handed neutrino superfield

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Abstract

In this work, we explore soft leptogenesis in the NMSSM framework extended by a right-handed neutrino superfield. We calculate the CP asymmetry, $\varepsilon$, and find it to be non-zero at tree level without using thermal effects for the final state particles. This is in contrast to soft leptogenesis in the MSSM extended by a right-handed neutrino superfield where thermal effects are essential. The difference arises due to the presence of a 3-body decay of the sneutrino in the NMSSM that violates lepton number at tree level. Apart from this, we also find that $\varepsilon \neq 0$ if the additional singlet scalar has a complex vacuum expectation value while all the other NMSSM parameters including the soft SUSY breaking ones relevant for CP asymmetry remain real. We estimate the order of magnitudes of these parameters to produce sufficient baryon asymmetry of the Universe.

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I. INTRODUCTION

It is well known that the observable Universe has an asymmetry between baryons and anti-baryons [1], often called the problem of baryogenesis. Over the years, many mechanisms have been proposed to create this baryon asymmetry. More recently, leptogenesis [2, 3] has become a highly favoured model for baryogenesis, specially because this mechanism is naturally linked to neutrino masses. Adding right-handed (RH) singlet heavy neutrinos to the standard model generates neutrino masses by the seesaw mechanism [4–9]. These RH neutrinos can also decay to produce a scalar and SM leptons and if the decay violates CP for example due to interference between tree-level and loop-level decays owing to complex couplings, a lepton asymmetry is generated. Since, in the standard model (SM), the $B - L$ symmetry is exact while the $B + L$ symmetry is broken by the electroweak (EW) sphaleron processes [10], these sphaleron processes can convert the generated lepton asymmetry to baryon asymmetry.

Soft leptogenesis [11–13] pertains to generating lepton asymmetry at the tree-level itself due to mixing between the particle and anti-particle states of the RH singlet sneutrino, $\tilde{N}$, because of the presence of soft SUSY breaking terms$^1$. In the most minimal soft leptogenesis setup using minimal supersymmetric standard model (MSSM) [25, 26] extended by one RH neutrino superfield, $\hat{N}$, a CP asymmetry in the RH sneutrino sector is created only when thermal masses for the final products are considered. The asymmetry is present because the thermal phase space factors are different for bosons and fermions. Another work featuring soft leptogenesis looks at CP violation not just due to mixing between particle and antiparticle initial states but in decays and in the interference of mixing and decay [27]. In Ref. [28], it is shown that considering most generic soft trilinear couplings and one loop self energy contributions for sneutrino decay it is possible to generate CP violation even without finite temperature effects within the same setup.

However, the MSSM suffers from the so-called $\mu$-problem [29] – there is no explanation to why the SUSY scale preserving $\mu$-term (a direct SUSY mass term for the Higgs fields) should be of the same order as the soft SUSY breaking terms. The most straightforward solution to the $\mu$-problem comes by promoting the $\mu$-parameter into a field whose vacuum expectation value (vev) is determined, like the other scalar field vevs, from the minimization of the

$^1$ Soft leptogenesis in different types of SUSY framework can be found in [14–24].
scalar potential along the new direction [30]. Naturally, it is expected to fall in the range of the other vevs, i.e., of order $O(M_{\text{SUSY}})$. The next-to-minimal supersymmetric standard model (NMSSM) (for review see [30, 31]) is the most simple and elegant model to solve this problem, where a singlet superfield $\hat{S}$ is introduced to the MSSM superfields which gets non-zero vev. The NMSSM can be extended by a set of RH neutrino superfields to generate masses for the SM light neutrinos by the type-I seesaw (see [32–34] for the MSSM extended by RH neutrino superfield). This has been explored earlier in Ref. [35]. This extension also keeps the R-parity conserved if the sneutrinos do not get vevs [35].

In this work, by using the NMSSM extended by the RH neutrino superfield, we present a soft leptogenesis scenario that creates a lepton asymmetry at the tree-level decay of the RH sneutrino without using thermal mass factors. The CP violation is achieved by the mixing between the particle and anti-particle states. We also show that it is possible to obtain non-zero CP asymmetry even when all the soft parameters are real. Since the soft terms are responsible for creating the CP asymmetry instead of needing flavour effects as in usual leptogenesis, using only one generation of the RH neutrino superfield is enough. Even so, the setup can be easily extended to get the experimentally observed SM neutrino mass hierarchies and their mixing angles pattern [36, 37].

The paper is organised in the following manner. We setup the model and segregate the parts required for soft leptogenesis in the next section (Sec. II). In the one following that, i.e., Sec. III, we calculate the CP asymmetry produced by decays of the various particle present in the model that contribute to non-zero CP asymmetry parameter $\varepsilon$ at the tree level. We talk about the decays of $\tilde{N}$ as well as the scalar $S$ in the model. In Sec. IV, we discuss the most crucial and important constraints and give a simple expression for $\varepsilon$ is given. In Sec. V, we give and discuss the results of our calculation. We find that for successfully generating the observed baryon asymmetry of the Universe, we need $\varepsilon \approx O(10^{-6})$. We also discuss what this could mean for various parameters of the model including the soft ones. We finally conclude in Sec. VI.

II. MODEL

In the NMSSM, an extra singlet superfield $\hat{S}$ is added to the MSSM Higgs sector [30]. Assuming explicitly $\mathbb{Z}_3$ symmetry, the superpotential for the NMSSM with a singlet RH
neutrino superfield $\hat{N}$ in terms of the new singlet superfield $\hat{S}$ and the MSSM doublet superfields $\hat{H}_u$ and $\hat{H}_d$ will be as follows [30]:

$$W = Y_E^i \hat{H}_d \hat{L}_i + Y_D^i \hat{H}_d \hat{D}_i + Y_U^i \hat{H}_u \hat{Q}_i + \lambda \hat{S} \hat{H}_u \hat{H}_d + \frac{\kappa}{3} \hat{S}^3 + Y_N^i \hat{N} \hat{H}_u \hat{L}_i + \lambda_N \hat{S} \hat{N} \hat{N},$$

(1)

where $\hat{L}_i$ and $\hat{Q}_i$ are the $SU(2)$ doublet superfields of leptons and quarks; $\hat{E}_i$ and $\hat{D}_i$ ($\hat{U}_i$) denote singlet down (up)-type quark superfields, respectively, and $Y$’s, $\lambda$’s and $\kappa$ are dimensionless couplings with generation indices $(i, j = 1, 2, 3)$. After the singlet $S$ obtains a vacuum expectation value (vev) $\langle S \rangle$, an effective $\mu$-term is generated: $\mu_{\text{eff}} = \lambda \langle S \rangle$, which solves the so-called $\mu$-problem [29]. The soft SUSY-breaking Lagrangian is given by

$$-\mathcal{L}_{\text{soft}} = -\mathcal{L}_{\text{soft}}^{\text{MSSM}}|_{B_{\mu} = 0} + \left( A_{\lambda_H} \lambda \hat{S} \hat{H}_u \hat{H}_d + A_{\kappa} \frac{\kappa}{3} \hat{S}^3 + A_{\kappa_N} Y_N^i \hat{N} \hat{H}_u \hat{L}_i + A_{\lambda_N} \lambda_N \hat{S} \hat{N} \hat{N} + \text{h.c.} \right) + m_S^2 |\hat{S}|^2 + M^2 |\hat{N}|^2,$$

(2)

where $\hat{L}_i$ and $\hat{N}$ are the scalar components of $\hat{L}_i$ and $\hat{N}$ superfields, respectively. CP is spontaneously violated when the scalars $\hat{H}_u, \hat{H}_d, S$ attain vevs with relative physical phases. The vev of the singlet $S$ is complex:

$$\langle S \rangle = v_S e^{i\delta}.$$

(3)

Since leptogenesis occurs above the electroweak (EW) phase transition, we do not give vevs to the two Higgs doublets. In this case, spontaneous CP violation can occur only when $\sin \delta \neq 0$.

### A. Terms relevant for soft leptogenesis

The terms from the superpotential required for leptogenesis via sneutrino decay are:

$$W \supset Y_N \hat{L} \hat{H} \hat{N} + \lambda_N \hat{S} \hat{N} \hat{N} + \frac{\kappa}{3} \hat{S}^3.$$

(4)

Here we consider $\lambda_N$, $\kappa$ to be all real and positive. We also remove the $i,j$ indices from the leptons and the $u$ index from the Higgs superfield for brevity. The scalar potential is obtained using:

$$V_S = \left| \frac{\partial W}{\partial S} \right|^2 + \left| \frac{\partial W}{\partial \hat{N}} \right|^2 + \left| \frac{\partial W}{\partial \hat{L}} \right|^2,$$

(5)
with
\[
\left| \frac{\partial W}{\partial S} \right|^2 \bigg| = \lambda_N^2 |\bar{N}|^4 + \kappa^2 |S|^4 + \lambda_N \kappa S^* \bar{N} N + \lambda_N \kappa \bar{N}^* \bar{N}^* S^2, \tag{6}
\]
\[
\left| \frac{\partial W}{\partial N} \right|^2 \bigg| = |Y_N|^2 |\bar{L}|^2 |H|^2 + 4\lambda_N^2 |\bar{N}|^2 |S|^2 + 2\lambda_N Y_N^* \bar{N} S \bar{L}^* H^* + 2\lambda_N Y_N \bar{N}^* S^* H \bar{L}, \tag{7}
\]
\[
\left| \frac{\partial W}{\partial L} \right|^2 \bigg| = |Y_N|^2 |H|^2 |\bar{N}|^2. \tag{8}
\]

The fermionic part of the Lagrangian is given by:
\[
\mathcal{L}_f = Y_N \bar{L} H \bar{N} + \lambda_N S N N + Y_N L H N. \tag{9}
\]

The soft SUSY-breaking Lagrangian terms that play a role in leptogenesis are:
\[
-\mathcal{L}_{\text{soft}} \supset \left( A_\kappa \frac{\kappa}{3} S^3 + A_N Y_N \bar{L} H \bar{N} + A_N \lambda N S N \bar{N} N + \text{h.c.} \right) + m_\sigma^2 |S|^2 + M^2 |\bar{N}|^2. \tag{10}
\]

The superpotential and the soft breaking terms combine to give the following interactions for $\bar{N}$ and $\sigma$ which could in principle contribute to soft leptogenesis due to mixing between the particle and anti-particle states through the soft terms:
\[
\mathcal{L}_{\text{int}} = \bar{N} \left( Y_N \bar{H} L + 2\lambda_N Y_N^* v_S e^{i\delta} H^* \bar{L}^* + 2\lambda_N Y_N^* \sigma \bar{H}^* \bar{L}^* + A_N Y_N \bar{H} \bar{L} \right)

+ \sigma \left( \lambda_N N N + A_N \lambda_N \bar{N} \bar{N} \right) + \text{h.c.}, \tag{11}
\]

where $\sigma = S - \langle S \rangle$.

### III. CP Asymmetry

Because of the soft terms as well as the vev of $S$, there is a mixing between particle and anti-particle states of the sneutrino and the singlet scalar $\sigma$ which is the dynamic part of $S$.

The squared mass matrices for the two of them are given by:
\[
\mathcal{M}_{\bar{N}}^2 = \begin{bmatrix}
M_1^2 & \lambda_N \kappa v_S^2 e^{2i\delta} + A_N \lambda_N v_S e^{-i\delta} \\
\lambda_N \kappa v_S^2 e^{-2i\delta} + A_N \lambda_N v_S e^{i\delta} & M_1^2
\end{bmatrix}, \tag{12}
\]
\[
\mathcal{M}_\sigma^2 = \begin{bmatrix}
m_\sigma^2 & 2\kappa^2 v_S^2 e^{2i\delta} + 2A_\kappa \kappa v_S e^{-i\delta} \\
2\kappa^2 v_S^2 e^{-2i\delta} + 2A_\kappa \kappa v_S e^{i\delta} & m_\sigma^2
\end{bmatrix}, \tag{13}
\]

where
\[
M_1^2 = M^2 + 4\lambda_N^2 v_S^2,
\]
\[
m_\sigma^2 = m_S^2 + 4\kappa^2 v_S^2. \tag{14}
\]
If $A_\lambda$ is real, the mass square eigenvalues of the sneutrino are:

$$M^2_{\pm} = M_1^2 \pm \sqrt{A^2_\lambda \lambda_N^2 v^2_S + \lambda_N^2 \kappa^2 v^4_S + 2 A_\lambda \lambda_N^2 v^2_S \kappa \cos(3\delta)},$$  \hfill (15)

which have the following eigenstates:

$$\tilde{N}_\pm = \frac{1}{\sqrt{2}} \left( \tilde{N} \pm \tilde{N}^* \right)$$  \hfill (16)

Similarly, one can write the mass square eigenvalues and eigenstates for the $\sigma - \sigma^*$ system. Because of mixing between the particle and anti-particle states of sneutrino and singlet scalar, these systems are similar to $K_0 - \bar{K}_0$ and $B_0 - \bar{B}_0$ systems [38]. The evolution of these systems in the non-relativistic limit are driven by the Hamiltonian $H$ defined as follows:

$$H = \mathcal{M} - \frac{i}{2} \Gamma,$$  \hfill (17)

where $\mathcal{M}$ is the mass matrix and $\Gamma$ is the decay rate matrix of the corresponding system.

Finally, the decay rates of the time evolved particle and anti-particle states of the $\tilde{N} - \tilde{N}^*$ and the $\sigma - \sigma^*$ system are calculated to get the final total CP asymmetry. Since both $\tilde{N}$ and $\sigma$ are massive and can decay into each other as final products as can be seen in Eq. (11), there are two possibilities for decay. The first is when $M_1 \gg m_\sigma$. In this case, $\tilde{N}$ decays to produce the CP asymmetry. Decays of $\sigma$ into a pair of right handed neutrino and sneutrinos is suppressed while if $m_\sigma \gg M_1$, the 3-body decay of $\tilde{N}$ into $\sigma, H^*$ and $\tilde{L}^*$ is suppressed. We consider these cases one-by-one.

**A. $\tilde{N}$ decays**

Irrespective of the mass of $\sigma$ relative to the mass of $\tilde{N}$, the sneutrino can decay into leptonic (sleptonic) and Higgs (Higgsino) final particles. The CP asymmetry generated from such a scenario was calculated in [11, 12]. However, if $M_1 \gg m_\sigma$, the 3-body decay channel opens up and it leads to interesting consequences for soft leptogenesis as we show below.

For the sneutrino system, upto leading order in the off-diagonal terms$^2$, the mass matrix

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$^2$ Considering $\kappa$ and $A_\lambda$ much smaller compared to $\lambda_N$ and $M_1$ respectively as will be justified later.
can be calculated from the squared mass matrix in Eq. (12):

$$\mathcal{M}_{\tilde{N}} = M_1 \begin{bmatrix}
1 & \frac{\lambda_N v^2 e^{2i\delta}}{2M_1^2} + \frac{A_N \lambda_N v e^{-i\delta}}{2M_1^2} \\
\frac{\lambda_N v^2 e^{-2i\delta}}{2M_1^2} + \frac{A_N \lambda_N v e^{i\delta}}{2M_1^2} & 1
\end{bmatrix}.$$  (18)

The decay rate matrix can be written from the Eq. (11). It contains both diagonal as well as off-diagonal terms because $\tilde{N}$ can decay into particle as well as anti-particle final states. There are both 2-body and 3-body decays of $\tilde{N}$ assuming that the mass of $\tilde{N}$ is much larger than the mass of $\sigma$.

$$\Gamma_{\tilde{N}} = \Gamma_1 \begin{bmatrix}
1 + \alpha + \frac{4\lambda_N^2 v^2}{M_1^2} + \frac{|A_N|^2}{M_1^2} & \frac{4\lambda_N v e^{-i\delta}A_N}{M_1^2} \\
\frac{4\lambda_N v e^{i\delta}A_N}{M_1^2} & 1 + \alpha + \frac{4\lambda_N^2 v^2}{M_1^2} + \frac{|A_N|^2}{M_1^2}
\end{bmatrix},$$  (19)

with

$$\Gamma_1 = \frac{|Y_N|^2 M_1}{8\pi},$$  (20)

$$\alpha = \frac{4\lambda_N^2}{\pi^2 M_1^2} \left[ M_1(M_1^2 + 4m_\sigma^2)^{1/2}/8 - \frac{1}{2} m_\sigma^2 \log \left( \frac{(M_1^2 + 4m_\sigma^2)^{1/2} + M_1}{2m_\sigma} \right) \right].$$  (21)

The solutions for the time evolution of $\tilde{N}$ and $\tilde{N}^*$ come from the Schrodinger like equation

$$\mathcal{H}\psi = i\frac{d\psi}{dt},$$  (22)

where $\psi = \{\tilde{N}, \tilde{N}^*\}^T$. The solutions are obtained as,

$$\tilde{N}(t) = e^{-i\alpha t} \left[ \tilde{N}_0 \cos \left( \frac{p}{q} bt \right) - i \tilde{N}_0^* \frac{q}{p} \sin \left( \frac{p}{q} bt \right) \right],$$  (23)

$$\tilde{N}^*(t) = e^{-i\alpha t} \left[ \tilde{N}_0^* \cos \left( \frac{p}{q} bt \right) - i \tilde{N}_0 \frac{p}{q} \sin \left( \frac{p}{q} bt \right) \right],$$  (24)

where $\tilde{N}_0, \tilde{N}_0^*$ are the field values at $t = 0$ and

$$a = (\mathcal{M}_{\tilde{N}})_{11} - i \frac{(\Gamma_{\tilde{N}})_{11}}{2} = (\mathcal{M}_{\tilde{N}})_{22} - i \frac{(\Gamma_{\tilde{N}})_{22}}{2},$$  (25)

$$b = (\mathcal{M}_{\tilde{N}})_{12} - i \frac{(\Gamma_{\tilde{N}})_{12}}{2},$$  (26)

$$\left( \frac{p}{q} \right)^2 = \frac{(\mathcal{M}_{\tilde{N}})^*_{12} - i \frac{(\Gamma_{\tilde{N}})^*_{12}}{2}}{(\mathcal{M}_{\tilde{N}})_{12} - i \frac{(\Gamma_{\tilde{N}})_{12}}{2}}.$$  (27)
Let’s define $\Delta M = M_+ - M_-$ and $\Delta \Gamma_\tilde{N} = \Gamma_+ - \Gamma_-$ and $Q = \frac{q}{p} b$. Then if $\Gamma_1^2 \ll M_1^2$ which happens when $Y_N \ll 1$ (typically of $\mathcal{O}(10^{-4})$ to satisfy neutrino mass bounds), we can write

$$2 \text{Re}(Q) \simeq \Delta M = \frac{\lambda_N v_S}{M_1} \left[ A_1^2 + \kappa^2 v_S^2 + 2 A_\lambda v_S \kappa \cos(3\delta) \right]^{1/2},$$

$$-4 \text{Im}(Q) \simeq \Delta \Gamma_\tilde{N} = \frac{2 Y_N^2 \lambda_N^2 v_S^2}{\pi^2 M_1^2 \Delta M} \left\{ k v_S (\cos(3\delta) \text{Im} A_N + \sin(3\delta) \text{Re} A_N) + A_\lambda \text{Re} A_N \right\}. \quad (29)$$

If $\Delta \Gamma_\tilde{N} \ll \Delta M$ as well, the argument of the trigonometric functions becomes $\Delta M t/2$ such that we can write:

$$\tilde{N}(t) = g_1 \tilde{N}_0 + \frac{q}{p} g_2 \tilde{N}_0^*, \quad (30)$$

$$\tilde{N}^*(t) = g_1 \tilde{N}_0^* + \frac{p}{q} g_2 \tilde{N}_0, \quad (31)$$

where

$$g_1 = e^{-i M_1 t} \exp \left[ -\frac{\Gamma_1}{2} \left( 1 + \alpha + \frac{4 \lambda_N^2 v_S^2}{M_1^2} + \frac{|A_N|^2}{M_1^2} \right) t \right] \cos \left[ \frac{\Delta M t}{2} \right], \quad (32)$$

$$g_2 = -ie^{-i M_1 t} \exp \left[ -\frac{\Gamma_1}{2} \left( 1 + \alpha + \frac{4 \lambda_N^2 v_S^2}{M_1^2} + \frac{|A_N|^2}{M_1^2} \right) t \right] \sin \left[ \frac{\Delta M t}{2} \right]. \quad (33)$$

The CP asymmetry factor $\varepsilon$ is defined as the ratio of the difference between the decay rates of $\tilde{N}$ and $\tilde{N}^*$ into final state particles with lepton number +1 and −1 to the sum of all the decay rates, i.e.,

$$\varepsilon = \frac{\sum_f \int_0^\infty dt \left[ \Gamma(\tilde{N}(t) \to f) + \Gamma(\tilde{N}^*(t) \to f) - \Gamma(\tilde{N}(t) \to \bar{f}) - \Gamma(\tilde{N}^*(t) \to \bar{f}) \right]}{\sum_f \int_0^\infty dt \left[ \Gamma(\tilde{N}(t) \to f) + \Gamma(\tilde{N}^*(t) \to f) + \Gamma(\tilde{N}(t) \to \bar{f}) + \Gamma(\tilde{N}^*(t) \to \bar{f}) \right]}, \quad (34)$$

where $f, \bar{f}$ are the final states with lepton number +1 and −1, respectively. This then gives us the following CP asymmetry parameter:

$$\varepsilon = \frac{\int_0^\infty dt \left[ |g_1|^2 |Y_N|^2 M_1 \right] \left[ 1 + \frac{|A_N|^2}{M_1^2} - \frac{4 \lambda_N^2 v_S^2}{M_1^2} - \frac{4 \lambda_N^2 \kappa}{\pi^2 M_1^2} \right]}{\int_0^\infty dt \left[ |Y_N|^2 M_1 \left[ 1 + \frac{|A_N|^2}{M_1^2} + \frac{4 \lambda_N^2 v_S^2}{M_1^2} + \frac{4 \lambda_N^2 \kappa}{\pi^2 M_1^2} \right] \right] \left[ |g_1|^2 + \frac{|g_2|^2}{2} \left( \left| \frac{q}{p} \right|^2 + \left| \frac{p}{q} \right|^2 \right) \right]}. \quad (35)$$

The decay rates at the tree level itself for final states with lepton numbers ±1 are different because the factor $\left| \frac{q}{p} \right| \neq 1$ as it is not a hermitian quantity. This requires non-zero off-diagonal terms to be present in the mass matrix as well as the decay rate matrix of the system with at least one of them being complex.
We also note that while the usual soft leptogenesis done in the MSSM [11, 12], to the leading order in soft terms, necessarily requires thermal phase space factors for the final state bosons \((c_B)\) and fermions \((c_F)\) to have \(\varepsilon \propto \Delta_{BF} = \frac{c_B - c_F}{c_B + c_F}\), we get an asymmetry even at zero temperature. This happens because of the presence of \(\sigma\) in the model which facilitates a 3-body decay which is not cancelled by the other terms. If we did not have this, at leading order in soft terms, \(1 - \frac{4\lambda^2 \nu^2}{M^2} = 0\) and there would be no asymmetry without thermal mass corrections.

The sum and difference of the ratios \(|q/p|^2\) and \(|p/q|^2\) can be written as:

\[
\begin{align*}
\left|\frac{q}{p}\right|^2 - \left|\frac{p}{q}\right|^2 &= -2 \left(\frac{y^2}{x^2 - y^2}\right)^{1/2}, \quad (36) \\
\left|\frac{q}{p}\right|^2 + \left|\frac{p}{q}\right|^2 &= 2 \left(\frac{x^2}{x^2 - y^2}\right)^{1/2}, \quad (37)
\end{align*}
\]

where

\[
\begin{align*}
x &= \frac{\lambda_N^2 \nu^4}{4} + \frac{2\lambda_N^2 \nu^2}{4} + \frac{4\Gamma_A^2 \lambda_N^2 \nu^3 |A_N|^2}{M^2} + \frac{\lambda_N^2 A_N \nu \cos(3\delta)}{2}, \quad (38) \\
y &= \frac{2\Gamma_A \lambda_N^2 \nu^3 \text{Im} A_N}{M^2} + \frac{4\lambda_N^2 \nu \kappa_S^3 \Gamma_A}{M^2} (\cos(3\delta) \text{Im} A_N + \sin(3\delta) \text{Re} A_N). \quad (39)
\end{align*}
\]

Keeping terms up to the leading order in \(\Gamma_1\), we find that the sum of the ratios is \(\simeq 2\) while the difference is twice the values of \(y/(x^2 - y^2)^{1/2}\) with,

\[
\left(\frac{y^2}{x^2 - y^2}\right)^{1/2} = \frac{|Y_N|^2 [A_N \text{Im} A_N + \kappa_S (\cos(3\delta) \text{Im} A_N + \sin(3\delta) \text{Re} A_N)]}{\pi [\nu^2 \kappa_S^2 + A_N^2 + 2A_N \kappa_S \cos(3\delta)]} \quad (40)
\]

such that the final CP asymmetry can be written as:

\[
\varepsilon = -\frac{\Delta M^2}{2(\Gamma^2 + \Delta M^2)} \left[1 + \frac{|A_N|^2}{M^2} - \frac{4\lambda_N^2 \nu^2}{M^2} - \frac{4\lambda_N^2 \nu \beta}{\pi^2 M^2}\right] \left(\frac{y^2}{x^2 - y^2}\right)^{1/2}. \quad (41)
\]

From Eqs. (40) and (41) it is evident that there is non-zero CP asymmetry even with real \(A_N\), provided \(\delta\) is sufficiently large.

Therefore it is possible to successfully generate non-zero lepton asymmetry from \(\tilde{N} - \tilde{N}^*\) system at the tree level without using thermal phase space factors for bosonic and fermionic final states. We discuss this more and give some numerical estimates in Sec. V for relevant parameters. For the moment, let’s consider the decays of \(\sigma\).
B. $\sigma$ decays

Unlike the $\tilde{N}$ decays where most of the decay products were massless, the final products of $\sigma$ decay are massive. This creates two possible decay modes of $\sigma$ according to the condition satisfied.

1. $\sigma$ decays to $NN$ and $\tilde{N}\tilde{N}$. This happens when $m_\sigma^2 > 4M_1^2$.

2. $\sigma$ decays only to $NN$. This happens when $16\lambda_N^2v_S^2 = 4m_N^2 < m_\sigma^2 < 4M_1^2$.

The mass matrix and the $\Gamma$ matrix of the $\sigma - \sigma^*$ system are respectively:

$$
M_\sigma = m_\sigma \begin{bmatrix}
1 & \frac{\kappa^2v_S^2e^{2i\delta} + A^*e^{-i\delta}}{m_\sigma^2} \\
\frac{\kappa^2v_S^2e^{-2i\delta} + A^*e^{i\delta}}{m_\sigma^2} & 1
\end{bmatrix},
$$

$$
\Gamma_\sigma = \Gamma_{\sigma,1} \begin{bmatrix}
\Theta & 0 \\
0 & \Theta
\end{bmatrix},
$$

where

$$
\Gamma_{\sigma,1} = \frac{\lambda_N^2}{32\pi m_\sigma^2},
$$

$$
\Theta = \begin{cases}
(m_\sigma^2 - 4m_N^2)^{3/2} + A^2(m_\sigma^2 - 4M_1^2)^{1/2} & \text{for case 1,} \\
(m_\sigma^2 - 4m_N^2)^{3/2} & \text{for case 2.}
\end{cases}
$$

A non-relativistic Hamiltonian can be defined following Eq. (17). Immediately it can be seen that because of the absence of an off-diagonal term in the decay rate matrix of $\sigma$, the ratio corresponding to $(p/q)^2$ of $\tilde{N}$ decay,

$$
\left(\frac{s}{r}\right)^2 = \frac{(M_\sigma)^*_1 - i\langle\Gamma_\sigma\rangle_{12}}{2(M_\sigma)_{12} - i\langle\Gamma_\sigma\rangle_{12}} = \frac{(M_\sigma)^*_1}{(M_\sigma)_{12}}.
$$

If we solve for the evolution of the $\sigma - \sigma^*$ system, the CP asymmetry parameter computed exactly analogously to the $\tilde{N} - \tilde{N}^*$ system will be zero because of exact cancellation between the ratios $|r/s|^2$ and $|s/r|^2$. Thus

$$
\varepsilon_\sigma = 0.
$$

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IV. GENERAL CONSTRAINTS

The CP asymmetry depends in this model depends on a lot of parameters. However, we can constrain some of them by various considerations. In deriving the following general constraints, we take $M_1 \gg M \Rightarrow M_1 \simeq 2\lambda_N v_S$, $M_1 \gg m_\sigma$ and $A_\lambda \simeq \kappa v_S$ for simplicity.

- The condition of out-of-equilibrium decays at $T = M_1$ is given by comparing the decay rate of the $\tilde{N}$ with the Hubble parameter at $T = M_1$:

$$\Gamma \lesssim H(T = M_1) = \sqrt{\frac{8\pi^3 g_s}{90} \frac{M_1^2}{m_{Pl}}}, \quad (48)$$

where $\Gamma$ is the diagonal component of Eq. (20). Substituting it in Eq. (48) we get

$$\frac{|A_N|^2}{M_1^2} \lesssim \frac{13\pi\sqrt{g} M_1}{Y_N^2 m_{Pl}} - 2 - \alpha. \quad (49)$$

For $M_1 \gg m_\sigma$, $\alpha \approx O(10^{-2})$ and we may write Eq. (49) as

$$\frac{|A_N|^2}{M_1^2} \lesssim \frac{13\pi\sqrt{g} M_1}{Y_N^2 m_{Pl}} - 2. \quad (50)$$

- The way we derived the CP asymmetry requires well separated states [12], i.e., $\Gamma \ll \Delta M$ as well as $\Delta \Gamma \ll \Delta M$ as stated before. This gives us two self-consistent limits:

$$\frac{|A_N|^2}{M_1^2} \ll \frac{8\pi\kappa v_S}{Y_N^2 M_1} - 2, \quad (51)$$

$$\cos(3\delta)\text{Im}A_N + (\sin(3\delta) + 1) \text{Re}A_N \ll \frac{2\pi^2}{Y_N^2} \kappa v_S. \quad (52)$$

- Neutrino mass upper limits ($m_\nu \lesssim 0.1$ eV [1, 39]) put constraints on the Yukawa coupling strength $Y_N$ and the mass of the right handed (s)neutrino.

$$\frac{Y_N^2}{\lambda_N v_S} \lesssim 6.6 \times 10^{-15} \text{ GeV}, \quad (53)$$

- Electric dipole moment calculations can constrain the CP violating phases that appear in the vevs of the two Higgs doublets and the scalar singlet $S$. In [40], they show that in principle $\delta_u$ (the phase in the vev of $H_u$ should we go below the EW scale) and $\delta$ could be large as long as the relative phase is kept small. Since we keep vev only for $S$, we keep $\delta$ small. For more details about the EDM constraints on the NMSSM, see [41].
A. A simpler form for $\epsilon$

We can write a simpler form for the CP asymmetry by using the approximations made and the general relationships between various parameters given above. The set of parameters governing $\epsilon$ are:

$$\{m_S, M, \kappa, \lambda_N, v_S, Y_N, A_\lambda, \text{Re}A_N, \text{Im}A_N, \delta\}.$$  

We choose the soft masses $M, m_S \sim \mathcal{O}(1) \text{ TeV}$ and the vev of $S$ to be $v_S \sim 10^7 \text{ GeV}$ with $\lambda_N \sim \mathcal{O}(1)$. This along with $\kappa \ll \lambda_N$ means both $M, m_S \ll 2\lambda_N v_S \simeq M_1$. In addition we also assume $A_\lambda \sim \kappa v_S$. We can put and upper limit on $Y$ from neutrino mass bounds i.e. $Y_N < \mathcal{O}(10^{-4})$ following Eq. (53). With these choices and approximations the form of $\epsilon$ can be simplified to

$$\epsilon \simeq - \left[ \frac{2\pi^2 |A_N|^2 - 4v_S^2}{2\pi^2 A_N^2 + 16\pi^2 v_S^2 + 4v_S^2} \right] \times \left[ \frac{|Y_N|^2}{8\pi \kappa v_S} \times \left[ \text{Im}A_N \{ 1 + \cos(3\delta) \} + \sin(3\delta) \text{Re}A_N \right] \right]$$

To have enough asymmetry ($\epsilon \simeq \mathcal{O}(10^{-6})$), it is evident that $\kappa \ll 1$ is required unless $A_N$ becomes very large. However large $A_N$ would violate the out-of-equilibrium decay condition in Eq. (50).

V. RESULTS & DISCUSSIONS

To obtain the baryon asymmetry of the Universe, $\eta_B$ we solve the simultaneous Boltzmann equations for the $\tilde{N}$ number density, $N_{\tilde{N}}$, and the $B - L$ number density, $N_{B - L}$ which are as follows [42–44]:

$$\frac{dN_{\tilde{N}}}{dz} = -K_{\tilde{N}} z (N_{\tilde{N}} - N_{\text{eq}}) \frac{\kappa_1(z)}{\kappa_1(z)},$$

$$\frac{dN_{B - L}}{dz} = -\epsilon K_{\tilde{N}} z (N_{\tilde{N}} - N_{\text{eq}}) \frac{\kappa_1(z)}{\kappa_2(z)} - \frac{1}{4} K_{\tilde{N}} z^3 \kappa_1(z) N_{B - L},$$

where $K_{\tilde{N}} = \frac{\Gamma}{H(z=1)}$ is the Hubble parameter at $z = 1$ with $z = \frac{M_1}{T}$, $N_{\text{eq}}$ is the equilibrium number density of $\tilde{N}$. They take the following forms:

$$H(z = 1) = \sqrt{\frac{8\pi^3 g_\ast}{90}} \frac{M_1^2}{m_{Pl}},$$

$$N_{\text{eq}} = \kappa_2(z) \frac{z^2}{2},$$
FIG. 1. The dependence of the baryon asymmetry on the CP asymmetry. In this figure, we keep $Y_N = 10^{-4.5}$, $\lambda_N = 0.9$, $M = 10^4$ GeV, $v_S = 10^{6.5}$ GeV, $|A_N| = 10^{6.5}$ GeV, $m_\sigma = 10^2$ GeV. The variation in $\varepsilon$ is brought on by not fixing $A_\lambda, \kappa, \delta$. We also take equilibrium initial condition for $\tilde{N}$ abundance. Starting with zero initial equilibrium for $\tilde{N}$ does not change the result.

with $m_{Pl} = 1.22 \times 10^{19}$ GeV being the Planck mass and $g_s$ is the number of relativistic degrees of freedom in NMSSM which we take $\approx 225$ except for the $\tilde{N}$ which is non-relativistic. In writing the Boltzmann equation for $B-L$ number density, we neglect the $\Delta L = 2$ scattering processes for washout and assume it is dominated mostly by inverse decays. The contribution to washout from the scattering processes is small because we are in the weak washout regime with $K_\tilde{N} \ll 1$. The final $B-L$ number density thus created, $N^f_{B-L}$ then converts to the baryon asymmetry by the sphaleron processes such that the ratio of the baryon number density to the photon number density, $\eta_B$ is:

$$\eta_B = \frac{3}{4} g_s^0 a_{sph} N^f_{B-L},$$

where $g_s \simeq g_s \simeq 225$, $g_s^0$ is the effective number of relativistic degrees of freedom at recombination and $a_{sph}$ is the sphaleron conversion factor. Since we will solve the Boltzmann equations numerically, we use the complete form of $\varepsilon$ given in Eq. (41).

In Fig. 1, we show the typical value of the CP asymmetry that satisfies the observed
baryon asymmetry of the Universe. It turns out that we need \( \varepsilon \simeq \mathcal{O}(10^{-6}) \) to get the correct observed baryon asymmetry while satisfying neutrino mass bounds. This value of \( \varepsilon \) is similar to the one obtained by other vanilla leptogenesis scenarios. The only difference is that usual leptogenesis occurs with decays at the loop level interfering with tree level decays due to complex Yukawa couplings the violate CP. In the soft leptogenesis, the Yukawa parameter could very well remain real as the source of CP asymmetry lies elsewhere – varying mixing between \( \tilde{N} \) and \( \tilde{N}^* \) states over time.

To get \( \varepsilon \simeq \mathcal{O}(10^{-6}) \) we fix the values of the following parameters in line with the earlier approximations and constraints:

\[
v_S = 10^7 \text{ GeV}, \quad \delta = 0.3, \quad M = 1 \text{ TeV}, \quad m_S = 1 \text{ TeV}, \quad \lambda_N = 1.
\] (61)

We also assume \( A_\lambda \simeq \kappa v_S \) for simplicity in finding the correct set of values for other parameters. Using these, we show the relation between \( A_N \) and \( \kappa \) for different values of \( Y_N \) in Figs. 2 and 3 as contour plots in \( \varepsilon \).

In Fig. 2 we take a complex \( A_N \) with \( \text{Im}A_N = \text{Re}A_N \) and vary \( \kappa \) and \( \text{Re}A_N \) for \( Y_N = 10^{-4.8} \) (left panel) and \( Y_N = 10^{-4.7} \) (right panel). We find that \( \kappa \lesssim 10^{-5} \) and \( \text{Re}A_N \gtrsim 10^7 \text{ GeV} \) are needed to generate \( \varepsilon \gtrsim 10^{-6} \). The cyan colored region is ruled out by the out-of-equilibrium

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**FIG. 2.** The contour plots of \( \varepsilon \) in \( \text{Re}A_N - \kappa \) plane considering \( A_N \) complex for (left panel) \( Y_N = 10^{-4.8} \) and (right panel) \( Y_N = 10^{-4.7} \).
FIG. 3. The contour plots of $\varepsilon$ in $\text{Re} A_N - \kappa$ plane considering $A_N$ real for (left panel) $Y_N = 10^{-4.8}$ and (right panel) $Y_N = 10^{-4.7}$.

condition. The exculded region gets enhanced for larger value of $Y_N$.

We keep $A_N$ real in Fig. 3 and vary $\kappa$ and $\text{Re} A_N$ for similar values of $Y_N$ as before. It’s clear that non zero asymmetry can be created even with real $A_N$ as long as $\delta \neq n\pi$. However compared to Fig. 2, we find that the bounds on $A_N$ and $\kappa$ in Fig. 3 are stronger. Both Figs. 2 and 3 satisfy neutrino mass bounds and the conditions of $\Gamma, \Delta \Gamma \ll \Delta M$.

VI. CONCLUSION

We have presented a new mechanism for soft leptogenesis in the context of the NMSSM with a singlet right handed neutrino superfield. Similar to soft leptogenesis in MSSM, we also generate CP asymmetry at the tree level owing to the CP violation occurring due to the difference between the mass and CP eigenstates similar to the $K^0 - \bar{K}^0$ or the $B^0 - \bar{B}^0$ systems. The difference lies in the fact that MSSM soft leptogenesis requires using thermal masses and phase space factors for boson and fermion final states without which there is no asymmetry. In the NMSSM where the singlet scalar $S$ takes a vev, an asymmetry can be generated even with zero temperature fields. Further if there is spontaneous CP violation in the system with $\sin \delta \neq 0$, lepton asymmetry can be created without using any other
complex parameter.

In the numerical analysis, we considered the mass scale of the RH sneutrino to be $10^7$ GeV. This is well below the limits imposed by the cosmological gravitino problem [45]. We found that to generate sufficient asymmetry, one of the soft trilinear coupling $A_N$ needs to be $\gtrsim 10^7$ GeV and $\kappa \lesssim \mathcal{O}(10^{-5})$. This also tells us that $A_\lambda \simeq \mathcal{O}(10^2)$ GeV. It would be interesting to explore the flavor effects in the present scenario that we leave for a future study.

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