PHYSICAL REVIEW D **67**, 123515 (2003)

Sleptogenesis

Rouzbeh Allahverdi
TRIUMF, 4004 Wesbrook Mall, Vancouver, British Columbia, Canada V6T 2A3

Bhaskar Dutta
Department of Physics, University of Regina, Regina, Saskatchewan, Canada S4S 0A2

Anupam Mazumdar
CHEP, McGill University, 3600 University Road, Montréal, Quebec, Canada H3A 2T8

(Received 27 January 2003; published 20 June 2003)

We propose that the observed baryon asymmetry of the Universe can naturally arise from a net asymmetry generated in the right-handed sneutrino sector at fairly low reheat temperatures. The initial asymmetry in the sneutrino sector is produced from the decay of the inflaton, and is subsequently transferred into the standard model (s)lepton doublet via three-body decay of the sneutrino. Our scenario relies on two main assumptions: a considerable branching ratio for the inflaton decay to the right-handed (s)neutrinos, and Majorana masses which are generated by the Higgs mechanism. The marked feature of this scenario is that the lepton asymmetry is decoupled from the neutrino Dirac Yukawa couplings. We exhibit that our scenario can be embedded within minimal models which seek the origin of a tiny mass for neutrinos.

DOI: 10.1103/PhysRevD.67.123515 PACS number(s): 98.80.Cq, 11.30.Fs, 12.60.Jv

I. INTRODUCTION

The consistency of the abundance of the light elements synthesized during the big bang nucleosynthesis (BBN) requires that the baryon asymmetry of the Universe (BAU) parametrized as $\eta_{\rm B}=(n_{\rm B}-n_{\bar{\rm B}})/s$, with s being the entropy density and n_B the number density of the baryons, be in the range $(0.3-0.9)\times 10^{-10}$ [1]. This asymmetry can be produced from a baryon symmetric universe provided three conditions are simultaneously met: B and/or L violation, C and CP violation, and departure from thermal equilibrium [2]. Any produced asymmetry will, however, be washed away by the standard model (SM) $(B+L)$-violating sphaleron transitions which are active from temperatures 10^{12} GeV down to 100 GeV [3], if $B-L=0$. Therefore an asymmetry in $B-L$, which is subsequently reprocessed by sphalerons, is generally sought in order to yield the net baryon asymmetry given by $B=a(B-L)$. Here a is a model-dependent parameter; in the case of the SM, $a=28/79$, while in the minimal supersymmetric standard model (MSSM), $a=32/92$ [4].

An attractive mechanism for producing $B-L$ asymmetry is from the decay of the heavy right-handed (RH) Majorana neutrinos [5]. Since the RH neutrinos are the SM singlets, a Majorana mass M_N, which violates the lepton number, is compatible with all of its symmetries, and hence can be arbitrarily large beyond the electroweak scale. This provides a natural explanation for the light neutrinos via the seesaw mechanism [6].

The lepton asymmetry can be generated from the interference between the tree-level and the one-loop diagrams in an out-of-equilibrium decay of the RH neutrinos in the early Universe, provided CP-violating phases exist in the neutrino Yukawa couplings. The asymmetry thus obtained will be partially converted into the baryon asymmetry via sphaleron effects. This is the standard lore for producing lepton asymmetry from on-shell RH neutrinos, commonly known as leptogenesis [5,7,8]. This can be accomplished in different ways.

In thermal leptogenesis scenario, RH neutrinos come into equilibrium with the primordial thermal bath through Yukawa interactions. The decay of the lightest RH neutrino easily satisfies the out-of-equilibrium condition by virtue of having a sufficiently small Yukawa coupling [8]. In a model-independent analysis in Ref. [9], the authors have parametrized thermal leptogenesis by four parameters; the CP asymmetry, the heavy RH neutrino mass, the effective light neutrino mass, and the quadratic mean of the light neutrino masses. The final result was that an acceptable lepton asymmetry could be generated with $T_{\rm R}\sim M_1=\mathcal{O}(10^{10})$ GeV, and $\Sigma_i m_{\nu,i}<\sqrt{3}$ eV.

This is marginally compatible with the upper bound on $T_{\rm R}$ allowed from thermal gravitino production in supersymmteric models [10]. Gravitinos with a mass $O(\text{TeV})$ decay long after nucleosynthesis and their decay products can change abundance of the light elements synthesized during BBN. For 100 GeV$\leqslant m_{3/2}\leqslant 1$ TeV, a successful nucleosynthesis requires $n_{3/2}/s\leqslant(10^{-14}-10^{-12})$, which translates into $T_{\rm R}\leqslant(10^{8}-10^{10})$ GeV [10,11]. The possibility of non-thermal gravitino production [12] does not give rise to any threat as described in [13,14]. It was also suggested that gravitinos can also be produced directly from the inflaton decay [15], and in the decay of heavy stable neutral particles [16], but the yielded bounds will not be severe.

An interesting alternative is non-thermal leptogenesis. This could happen in many ways. The simplest possibility is to produce on-shell RH neutrinos, with a considerable branching ratio, in inflaton decay [17]. It is also possible to produce heavy RH neutrinos (even heavier than the inflaton) via preheating [18]. However, non-thermal leptogenesis is rather model dependent. For example, just fermionic preheating is plagued by the fact that the running coupling of the inflaton to the fermions can easily give rise to correction in the inflaton mass, which leads to the instabilities during the inflaton oscillations, described as in Ref. [19]. The inflaton condensate fragments as a result of that and forms interesting solitons.

In supersymmetric models there are additional options as one can also excite sneutrinos [20]. In fact sneutrinos are produced more abundantly than neutrinos during preheating [21]. Another possibility is creating a condensate of sneutrinos which yields the right asymmetry through its decay [22], or via Affleck-Dine mechanism [23].

Recently it has been noticed that successful leptogenesis does not require on-shell RH (s)neutrinos [24,25]. A minimal model was proposed in Ref. [25], where the lepton asymmetry is directly generated from the inflaton decay into the Higgs boson and leptons via off-shell RH (s)neutrinos. This model naturally results in a sufficiently low reheat temperature, and yields desirable baryon asymmetry for a rather wide range of inflationary scale, neither invoking preheating in a particular model nor any unnaturally suppressed couplings.

In this paper we propose a completely new scenario for leptogenesis, called sleptogenesis.[1] We show that an asymmetry between sneutrinos and anti-sneutrinos can be generated, through a phase mismatch between the inflaton coupling to the RH (s)neutrinos and the Majorana masses, in inflaton decay. Note that the RH neutrino and antineutrino are indistinguishable due to the Majorana nature of neutrinos. After the (s)neutrinos decay, the SM (s)leptons carry the produced asymmetry which will be partially reprocessed to the baryon asymmetry. This scenario can emerge quite naturally provided the branching ratio for the inflaton decay to the RH (s)neutrinos is considerable, and there exists new Higgs field(s) generating the Majorana masses. The first assumption is rather common in non-thermal scenarios of leptogenesis, while the latter is necessary in models where the RH (s)neutrinos are gauge non-singlet under some new physics. The main feature of our scenario is replacing the dependence of the generated asymmetry on the neutrino Dirac Yukawa couplings with that on the Majorana Yukawa couplings. As a consequence, it is in principle possible to accommodate low-scale leptogenesis [27] with an appropriate choice of model parameters. The minimal extension of MSSM that can accommodate the above mentioned Majorana sector is $SU(3)_c\times SU(2)_L\times U(1)_{I_{3R}}\times U(1)_{B-L}$. In this scenario, the RH (s)neutrino masses arise at the scale where $U(1)_{I_{3R}}\times U(1)_{B-L}\rightarrow U(1)_Y$. The branching ratios of lepton flavor violating decay modes, e.g. $\tau\rightarrow\mu\gamma$, $\mu\rightarrow e\gamma$, will be able to discern these models in the near future.

II. THE SCENARIO

We begin by considering a simple model in a supersymmetric setup. The relevant part of the superpotential is given by

$$W\supset\frac{1}{2}m_\phi\Phi^2+\frac{1}{2}m_\sigma\Sigma^2+\frac{1}{2}\mathbf{y}\Phi\mathbf{N}^2+\frac{1}{2}\mathbf{g}\Sigma\mathbf{N}^2+\mathbf{h}\mathbf{N}\mathbf{H}_u\mathbf{L}+h_t\mathbf{H}_u\mathbf{Q}_3\mathbf{t}^c. \tag{1}$$

Here Φ is a gauge singlet superfield which comprises the inflaton ϕ and its superpartner (inflatino) with mass m_ϕ, and $\mathbf{N}$ is the superfield comprising the RH neutrino N and sneutrino $\tilde{N}$. While Σ comprises the scalar field σ which generates Majorana mass for N through its vacuum expectation value (VEV), denoted as σ_0, and its fermionic partner $\tilde{\sigma}$. As we will describe later, in realistic particle physics models $\mathbf{N}$ and Σ are charged under some gauge group (as a matter of fact, one needs to introduce another superfield $\overline{\Sigma}$ for anomaly cancellation). Since the inflaton is assumed to be a gauge singlet, its coupling to RH (s)neutrinos actually arises at the non-renormalizable level, and hence is small $\mathbf{y}\sim\mathcal{O}(m_\phi/M_{\rm P})$ (we use the reduced Planck mass $M_{\rm P}\sim 2.4\times 10^{18}$ GeV). This coupling will be responsible for decay of the inflaton to N and $\tilde{N}$ and, subsequently, reheating the Universe.

Finally, $\mathbf{H}_u$, $\mathbf{L}$, $\mathbf{Q}_3$, and $\mathbf{t}^c$ are the multiplets containing the Higgs boson which gives mass to the top quark, the left-handed lepton doublet, the third generation quark doublet and the RH top anti-quark, along with their superpartners, respectively. We have omitted all indices on $\mathbf{N}$, and lepton doublets. Note that $\mathbf{y}$ and $\mathbf{g}$ are symmetric matrices. For simplicity, we assume that they can be diagonalized in the same basis, and hence only their diagonal elements y_i and g_i are relevant.

We also assume that $m_\sigma\geqslant 10m_\phi$. This implies that the dynamics of σ is frozen during and after inflation, and hence ensures a simpler dynamics by virtue that all of the energy density is carried by ϕ. However, the mass of the RH (s)neutrinos M_i (at least one of them) is taken to be smaller than m_ϕ, so that the inflaton decay to N_i and $\tilde{N}_i$ will reheat the Universe.[2]

An important point is that the interference between the tree-level and one-loop contributions to the decay process $\phi\rightarrow\tilde{N}\tilde{N}$ results in an excess, or deficit, of $\tilde{N}$ over $\bar{\tilde{N}}$, provided a relative phase exists between g_i and y_i. This happens in exactly the same fashion as N decay generates a lepton asymmetry in the standard leptogenesis scenario [5].

Note that it is meaningless to talk of any asymmetry between N and $\overline{N}$, since there is no distinction between particle

[1]Baryogenesis with scalar fields has also been studied in Ref. [26], though in a different context.

[2]Note that we have neglected another coupling of the form $f\Phi\Sigma^2$, even though it can arise at the renormalizable level in realistic models, and hence need not be very small. The reason is that in the limit $m_\sigma\gg m_\phi$, such a coupling can only affect the inflaton decay by inducing $\phi\rightarrow\tilde{N}\tilde{N}\tilde{N}\tilde{N}$, via off-shell σ and $\tilde{\sigma}$, and $\phi\rightarrow\tilde{N}\tilde{N}$ decay modes at the tree-level and one-loop level, respectively. The effective coupling for these modes will be $f(gm_\phi/M)^2$ and fy, respectivley, and, moreover, their decay rate is suppressed by a four-body phase space factor and a one-loop factor, respectively. Thus the inflaton predominantly decays via coupling y, and a coupling between Φ and Σ will have no bearing on our results.

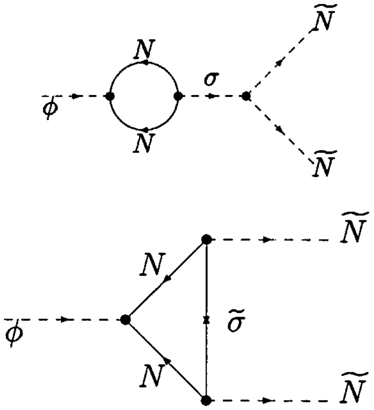

FIG. 1. One-loop self-energy and vertex diagrams resulting in an asymmetry between $\tilde{N}$ and $\bar{\tilde{N}}$.

and anti-particle for a Majorana fermion. To put it another way, the mass term $M_N NN$, which violates the lepton number, makes particle and anti-particle indistinguishable. On the other hand, the supersymmteric mass term $M_N^2|\tilde{N}|^2$ for the sneutrino does not violate the lepton number.

In most of the realistic models of inflation only the real component of the inflaton has a VEV. Then it can be shown from Eq. (1) that $\phi \to NN$ and $\phi \to \tilde{N}\tilde{N}$ decays occur at the same rate, and the total decay rate is given by

$$\Gamma_{\rm d} \simeq \frac{1}{8\pi} \sum_i y_i^2 m_\phi . \tag{2}$$

Note that $\Delta L = 2$ in $\phi \to \tilde{N}\tilde{N}$ decay. By taking into account the one-loop self-energy and vertex diagrams, shown in Fig. 1, we find that[3]

$$\frac{n_{\tilde{N}_i} - n_{\bar{\tilde{N}}_i}}{n_\phi} = -\frac{1}{8\pi} \frac{\mathrm{Im}[(\mathbf{y}\mathbf{g}^\dagger)_{ii}]^2}{\sum_i (\mathbf{y}\mathbf{y}^\dagger)_{ii}} f\left(\frac{m_\sigma^2}{m_\phi^2}\right), \tag{3}$$

where

$$f(x) = \sqrt{\frac{x}{2}}\left[\frac{2}{x-1} + \ln\left(1 + \frac{1}{x}\right)\right]. \tag{4}$$

These diagrams are similar to those in leptogenesis via $\tilde{N}$ decay [20] (with proper replacements). The expression for the asymmetry parameter therefore has exactly the same structure as in the standard leptogenesis [28]. There are slight differences though between the two cases. Here only half of the inflatons decay to RH sneutrinos, and ϕ decay to N does not lead to any asymmetry. On the other hand, the lepton number is violated by two units in $\phi \to \tilde{N}\tilde{N}$ decay. Finally, a factor of 1/2 arises in our case since identical particles appear in the loop. Note that in the limit $m_\sigma \geqslant 10 m_\phi$, we simply have $f \simeq 3m_\phi/2m_\sigma$.[4]

The created asymmetry is then transferred into the SM (s)leptons via $\tilde{N}_i$ decay. There are two two-body decay channels read from Eq. (1): $\tilde{N}_i \to \overline{L_i}\bar{\tilde{H}}$ and $\tilde{N}_i \to \tilde{L}_i H$, which have the same rate. Here h_i denotes diagonal elements of the neutrino Yukawa matrix $\mathbf{h}$ and, for simplicity, we assume that non-diagonal elements can be neglected. Since the two-body decays produce the same number of anti-leptons as leptons, no net lepton asymmetry will be yielded.

However, there exists a term $h_i h_t \tilde{N}_i \tilde{L} \bar{\tilde{Q}}_3 \bar{\tilde{t}}^c$ in the scalar potential which results in the three-body decay $\tilde{N}_i \to \bar{\tilde{L}} \tilde{Q}_3 \tilde{t}^c$. This channel is responsible for transferring the asymmetry into the SM (s)leptons, though with suppression by a factor $\simeq 3/32\pi^2$ (note that $h_t \approx 1$). The $1/32\pi^2$ is the ratio of phase space factors for three-body decay to the total decay rate, and note that $\tilde{N}$ decays to all three colors of squarks. In addition, we also have the usual dilution due to the entropy release from reheating by a factor of $T_{\rm R}/m_\phi$, where $T_{\rm R}$ denotes the reheat temperature. A thermal bath of the SM particles (and their superpartners) is typically formed right after $\tilde{N}$ and N decay (for details on thermalization, see Ref. [30]), and hence $T_{\rm R}$ is determined by the details of these decays.

Here we assume that all $\tilde{N}_i$ (and N_i) decay very rapidly right after they have been produced. This will simplify the calculations while preserving the essence of our scenario. It will be the case if $\Gamma_i \geqslant \Gamma_{\rm d}$, where Γ_i is the decay rate of $\tilde{N}_i$ (and, by virtue of supersymmetry, N_i). The (s)neutrinos, with mass M_i, initially having an energy $\simeq m_\phi/2$, and hence their decay rate (at the time of production) is given by

$$\Gamma_i \simeq \frac{h_i^2 M_i^2}{2\pi m_\phi}. \tag{5}$$

Note that the decay rate at the $\tilde{N}_i$ rest frame is $h_i^2 M_i/4$, and the time-dilation factor will be $2m_\phi/M_i$.

The requirement that the (s)neutrinos decay when $H \simeq \Gamma_{\rm d}$ translates into the condition $4h_i^2 M_i^2 \geqslant y^2 m_\phi^2$, where $y^2 = \Sigma_i y_i^2$. In the minimal seesaw model the limit on the light neutrino masses, with the current cosmological and laboratory bounds on the absolute neutrino masses taken into account, translates to

$$\frac{h_i^2 \langle H_u^0 \rangle^2}{M_i} \leqslant 10^{-9}\ \mathrm{GeV}, \tag{6}$$

[3]There are also contributions from supersymmetry breaking terms to these diagrams which will be suppressed as $m_{3/2}/m_\sigma$.

[4]In the limit $m_\phi = m_\sigma$ the perturbative results in Eqs. (3), (4) break down. In this case one has to actually take into account the finite decay width of ϕ and σ. This has been done for the standard lepotogenesis with degenerate Majorana (s)neutrinos, and it is shown that no asymmetry will be yielded, as expected, in the $x = 1$ limit [29].

where $\langle H_u^0\rangle \simeq 174$ GeV is the Higgs boson VEV. Since $M_i < m_\phi$, the instant (s)neutrino decay requires that $y^2 < 10^{-14}(m_\phi/1\ \mathrm{GeV})$. This results in a tiny y, which also fulfills the requirement from the model building point of view. A small coupling y also ensures a sufficiently low T_R.

After setting all the pieces together, including the reprocessing by sphalerons and dilution from reheating, we obtain

$$\eta_\mathrm{B} \simeq \frac{9}{64\pi^2}\frac{1}{8\pi}\frac{\sum_i y_i^2 g_i^2}{y^2}\frac{T_\mathrm{R}}{m_\sigma}, \tag{7}$$

where

$$T_\mathrm{R} \simeq \frac{g_*^{1/4}}{3}(y^2 M_\mathrm{P} m_\phi)^{1/2}. \tag{8}$$

Here g_* is the number of relativistic degrees of freedom ($g_* \simeq 200$ in the MSSM when $T_\mathrm{R} > 1$ TeV). Note that $n_\phi \simeq g_* T_\mathrm{R}^4/3m_\phi$, while $s \simeq g_* T_\mathrm{R}^3/\pi^2$.

Let us denote $\tilde{N}_1$ as the sneutrino which makes the largest contribution to the asymmetry. Then Eq. (7) implies that it has the largest combination yg, but not necessarily the largest y or g. Note that the inflaton mainly decays into the (s)neutrino with the largest y, while the heaviest (s)neutrino has the largest coupling g [see Eq. (1)]. The maximum asymmetry is yielded when $y_1 > y_2, y_3$. For $y^2 \simeq y_1^2$ the expression in Eq. (7) is further simplified to

$$\eta_\mathrm{B} \simeq \frac{g_1^2}{2^9\pi}\frac{T_\mathrm{R}}{m_\sigma}. \tag{9}$$

Therefore a successful leptogenesis requires that

$$g_1^2 \frac{T_\mathrm{R}}{m_\sigma} \gtrsim 5\times 10^{-8}. \tag{10}$$

A couple of important comments are in order now. The preservation of the lepton number by the sneutrino mass term has been a key point in our scenario. This is true for the sneutrino supersymmetric mass derived from the superpotential. However, supersymmetry must be broken in any realistic model and this inevitably introduces soft breaking terms. The soft breaking mass term $m_{3/2}^2|\tilde{N}|^2$, with $m_{3/2}$ being the gravitino mass, also preserves the lepton number. On the other hand, the A-term associated with the Majorana mass term, which has the form $a m_{3/2} M_N \tilde{N}\tilde{N} + \mathrm{H.c.}$, breaks the lepton number in the sneutrino sector. This term will cause an oscillation between the sneutrino and anti-sneutrino, similar to the neutrino flavor oscillations, with a frequency $a m_{3/2}$. In consequence, any asymmetry between $\tilde{N}$ and $\bar{\tilde{N}}$ only survives for a time $\leqslant (a m_{3/2})^{-1}$, while being washed out by $\tilde{N} - \bar{\tilde{N}}$ oscillations at longer time scales. Therefore the success of our proposed scenario requires that $\tilde{N}_1$ decay early enough, i.e. $\Gamma_1 \geqslant a m_{3/2}$.

The value of $m_{3/2}$ depends on the mechanism for communicating supersymmetry breaking to the observable sector. In gravity-mediated models $m_{3/2} \simeq 100\ \mathrm{GeV} - 1$ TeV, while in gauge-mediated models substantially smaller values $m_{3/2} \simeq 1$ KeV are possible. The situation then depends on the exact value of a, which is determined by the structure of the Kähler potential. For minimal Kähler terms one typically has $a \simeq \mathcal{O}(1)$, while $a \approx 0$ can be obtained in non-minimal cases. Let us focus on the former case, as it will clearly result in a more stringent bound. Then it is required that

$$\frac{h_1^2}{8\pi} M_1 \geqslant 10^2 \ \ (10^{-6}) \ \ \mathrm{GeV}, \tag{11}$$

in gravity (gauge)-mediated models in order to preserve the lepton asymmetry. By taking into account the see-saw constraint in Eq. (6), we obtain the absolute lower bound

$$M_1 \geqslant 10^{8.5} \ \ (10^{4.5}) \ \ \mathrm{GeV}, \tag{12}$$

on the mass of the RH neutrino with the largest contribution to the asymmetry. Note that the above bound is only meant for the minimal Kähler structure and can be significantly weakened for non-minimal kinetic terms.

III. WASHOUT OF THE GENERATED ASYMMETRY

We now turn our attention to various interactions which can wash out the produced asymmetry. First, let us briefly recount the thermal history of the Universe in our scenario. The inflaton mainly decays into the $\mathbf{N}_1$ multiplet when $H \simeq \Gamma_\mathrm{d}$, and a lepton asymmetry is generated in the decay to the sneutrino component $\tilde{N}_1$. Then $\tilde{N}_1$, as well as other (s)neutrinos, decays promptly and we obtain a thermal bath consisting of the SM degrees of freedom (and their superpartners) with temperature T_R estimated in Eq. (8).

The first lepton-number violating interaction is the N and $\tilde{N}$-mediated scattering of leptons and Higgs bosons (also their superpartners) in a thermal bath. These scatterings have been considered in detail in the standard leptogenesis scenario [31,7,8]. As an illustration, a sample scattering of this type will be inefficient only if

$$\Gamma_{\not L} \simeq \frac{h^4}{16\pi^3}\frac{T_\mathrm{R}^3}{M_N^2} < g_*^{1/2}\frac{T_\mathrm{R}^2}{M_P}. \tag{13}$$

Note that there exists a large number of such scattering, especially in the MSSM [8].

By using the relationship in Eq. (6), we obtain the constraint on reheat temperature which will avoid erasure of the lepton asymmetry. This bound turns out to be smaller than the gravitino overproduction bound $T_\mathrm{R} < 10^{10}$ GeV.

There are also other lepton number violating interactions, namely, the $\tilde{\sigma}$ and σ-mediated $\tilde{N}_1\tilde{N}_1$ and $\tilde{N}_1 N_1$ scatterings, shown in Fig. 2.[5] These processes can erase the lepton asymmetry carried by $\tilde{N}_1$ before it decays, provided they occur at

[5] Note that $\tilde{\phi}$ and ϕ-mediated scatterings can be neglected due to the smallness of the inflaton coupling to $\tilde{N}_1$ and N_1.

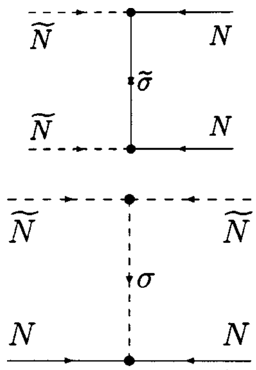

FIG. 2. Processes violating the lepton number in the sneutrino sector.

a higher rate. Note that the number density of N_1 and $\tilde{N}_1$ is $\simeq g_* T_R^4/3m_\phi$. This will result in

$$\Gamma_{\tilde{N}_1 N_1} \simeq \frac{Cg_1^4}{24\pi^3}\frac{g_* T_R^4}{m_\phi m_\sigma^2}, \tag{14}$$

and

$$\Gamma_{\tilde{N}_1 \tilde{N}_1} \simeq \frac{Cg_1^4}{12\pi^3}\frac{g_* T_R^4}{m_\phi m_\sigma^2}, \tag{15}$$

where C is a multiplicity factor representing different contribuitions to the same process, and $C/\pi^2 \sim \mathcal{O}(1)$. Also recall the decay rate $\Gamma_1 \simeq h_1^2 M_1^2/4\pi m_\phi$ for $\tilde{N}_1$. With the help of Eqs. (6), (10) we find that these processes will be inefficient, provided

$$\left(\frac{T_R}{1\ \text{GeV}}\right)^2 < 10^{-2}\left(\frac{M_1}{1\ \text{GeV}}\right)^3. \tag{16}$$

Note that $T_R < M_1$ for a perturbative decay of N_1 (for details see Ref. [30]). Therefore this bound is easily satisfied as long as $M_1 > 100$ GeV. In conclusion, the only non-trivial constraint in our scenario will be that of generating sufficient asymmetry, given in Eq. (10).

So far we have only considered the $\tilde{N}_1\tilde{N}_1$ and $\tilde{N}_1 N_1$ scatterings. On the other hand, $\tilde{N}_1\tilde{N}_1 \to N_i N_i$ and $\tilde{N}_1 N_1 \to \tilde{N}_i N_i$ annihilaitions can also happen through diagrams in Fig. 2. The rate for such processes is $\propto g_1^2 g_i^2$, which will be larger than the one considered above, provided N_1 is not the heaviest RH neutrino (note that $M_i \propto g_i$). However, as we shall see shortly, successful baryogenesis requires that M_1 not be much smaller than m_ϕ. This implies that M_1 is not very different from the largest $M_i < m_\phi$, and hence the rate for various processes represented by diagrams in Fig. 2 are in general comparable. Moreover, Eq. (14) will indeed give the largest rate if N_1 is the heaviest RH neutrino.

IV. MODEL PARAMETERS

We can now estimate the range of parameters within which our scenario can accommodate a successful baryogenesis. As an example, $m_\sigma \gtrsim 10 m_\phi$, which guarantees that σ does not play any dynamical role in the post-inflationary era, while from Eq. (16), $M_1 \gtrsim 10 T_R$ guarantees the survival of generated asymmetry. Then the observed baryon asymmetry can be obtained provided

$$g_1^2 \frac{M_1}{m_\phi} \gtrsim 5\times 10^{-6}. \tag{17}$$

For $g_1 \gtrsim 10^{-2}$, this would require that M_1 be (at least) an order of magnitude smaller than m_ϕ. This is at par with the standard non-thermal leptogenesis where (s)neutrinos are produced perturbatively. Note that a smaller M_1/m_ϕ is allowed as g_1 increases.

It is important to notice that, contrary to the standard leptogenesis scenario, sufficient asymmetry can be obtained with much smaller values of M_1. In fact, it is evident from Eq. (10) that η_B only depends on the ratio M_1/m_ϕ. Therefore, as advertised earlier, our scenario can accommodate low scale leptogenesis without making unnatural assumptions (e.g. having highly degenerate Majorana neutrinos, Ref. [27]). This is a consequence of generating the lepton asymmetry directly in the inflaton decay, and hence decoupling it from the neutrino Dirac Yukawas. One should, nevertheless, keep in mind the lower bound on M_1, from Eq. (12), which arises for the minimal Kähler potential. However, this has an entirely different origin, namely to avoid the earsure of the asymmetry by $\tilde{N}-\tilde{\bar{N}}$ oscillations induced by soft supersymmetry breaking terms. Moreover, it can be substantially weakened for a non-minimal Kähler structure.

V. EMBEDDING IN REALISTIC MODELS

The RH neutrino sector in Eq. (1) can be naturally added by extending the MSSM to incorporate a gauged $U(1)_{B-L}$ symmetry. Three fermions, with the same quantum number as the RH neutrinos, will then be required for gauge anomaly cancellation. The RH neutrinos obtain Majorana mass through the scalar component of the Σ superfield (with a $B-L$ charge of 2), which spontaneously breaks $U(1)_{B-L}$ symmetry. The present neutrino oscillation data indicates the scale of symmetry breaking v_{B-L} be somewhere around $10^{12}-10^{15}$ GeV. The presence of heavy RH neutrinos will ensure the light SM neutrino masses via the seesaw mechanism [6].

Note that the inflaton is considered to be a gauge singlet, and does not share any charge with other multiplets in Eq. (1). Thus its coupling to the RH neutrino sector is determined by non-renormalizable terms which, after symmetry breaking, result in $y \sim \mathcal{O}(v_{B-L}/M_P)$.

The simplest extension of the electroweak sector has the gauge group $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$, with the fermion quantum numbers assigned as follows: $\mathbf{Q}(2,0,+\frac{1}{3})$; $\mathbf{L}(2,0,-1)$; $\mathbf{u}^c(1,-\frac{1}{2},-\frac{1}{3})$; $\mathbf{d}^c(1,+\frac{1}{2},-\frac{1}{3})$; $\mathbf{e}^c(1,+\frac{1}{2},+1)$; $\mathbf{N}(1,-\frac{1}{2},+1)$. As mentioned earlier, three N are required from anomaly cancellations conditions. The Higgs fields have the assignment $\mathbf{H}_u(2,+\frac{1}{2},0)$; $\mathbf{H}_d(2,-\frac{1}{2},0)$; $\Sigma(1,+1,-2)$, $\bar{\Sigma}(1,-1,+2)$. Note that with the above charge assignments, two superfields Σ and $\bar{\Sigma}$ are required for anomaly cancellation. The mixings and mass differences among different neutrino flavors as observed in different experiments can be generated in this model via flavor violating Majorana couplings [32]. Indeed it is possible to find good fits of the experimental data with Majorana masses $>10^8$ GeV [32]. The branching ratios of lepton flavor violating decay modes, e.g. $\tau\to\mu\gamma$, $\mu\to e\gamma$ can distinguish these models.

VI. CONCLUSION

In this paper we have proposed a leptogenesis scenario where the lepton asymmetry is created in the RH sneutrino sector at relatively low reheat temperatures. This happens via a phase mismatch between the Majorana masses and the coupling of the RH (s)neutrinos to a gauge singlet inflaton. The prompt decay of the sneutrinos then transfers the lepton asymmetry to the SM lepton sector. The realization of this scenario requires a considerable branching ratio for the inflaton decay to (at least one of) the RH (s)neutrinos, and new Higgs boson field(s) whose VEV is responsible for generating the Majorana masses. The first requirement is a typical ingredient of non-thermal leptogenesis scenarios. The second one will be a necessary part of model building when the RH (s)neutrinos have gauge quantum charges under some new physics, e.g. models with a gauged $U(1)_{B-L}$ symmetry. The mixings and mass differences among different neutrino flavors as observed are generated in this model via flavor violating Majorana couplings. The remarkable difference from the standard leptogenesis is that here the asymmetry depends on the neutrino Majorana Yukawa couplings rather the Dirac Yukawas. There exists another source for the washout of the asymmetry in this scenario, in addition to the usual lepton number violating scatterings of leptons and Higgs bosons, namely the scattering of RH sneutrinos off each other or RH neutinos. We saw that for reheat temperatures compatible with the limit from thermal gravitino production, the washout processes do not lead to any meaningful constraints on the model parameters.

The maximum asymmetry is yielded when heavier (s)neutrinos have larger couplings to the inflaton. In this case the lepton asymmetry is mainly created in inflaton decay to the heaviset RH sneutrino $\tilde{N}_1$ with mass M_1. An acceptable baryon asymmetry can then be obtained for moderate Majorana Yukawa couplings $g_1 \gtrsim 10^{-2}$, and $m_\phi \gtrsim 10 M_1$.

One important point is that the low-energy supersymmetry breaking induces the $\tilde{N}-\bar{\tilde{N}}$ oscillations and, in consequence, erases the initial lepton asymmetry. This demands that the decay rate of RH sneutrinos must be larger than the frequency of such oscillations. The latter quantity depends on the form of the Kähler potential, as well as the mechanism for mediation of supersymmetry breaking. For minimal Kähler structure we require $M_1 > 10^{8.5}$ GeV in gravity-mediated models, and $M_1 > 10^{4.5}$ GeV in gauge-mediated models. These bounds can be substantially weakened for non-minimal cases. There will be no other constraints on M_1 besides this, and hence low scale leptogenesis can, in principle, be accommodated with a proper choice of the inflationary model.

ACKNOWLEDGMENTS

A.M. thanks CITA for financial support. He also acknowledges the hospitality of the Department of Physics, University of Regina where part of this work has been carried out. The research of R.A. and B.D. is supported by the National Sciences and Engineering Research Council of Canada.

[1] For a review, see K.A. Olive, G. Steigman, and T.P. Walker, Phys. Rep. **333**, 389 (2000).

[2] A.D. Sakharov, Pis'ma Zh. Éksp. Teor. Fiz. **5**, 32 (1967) [JETP Lett. **5**, 24 (1967)].

[3] V.A. Kuzmin, V.A. Rubakov, and M.E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985).

[4] S.Y. Khlebnikov and M.E. Shaposhnikov, Nucl. Phys. **B308**, 885 (1988).

[5] M. Fukugita and T. Yanagida, Phys. Lett. B **174**, 45 (1986).

[6] M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by P. van Nieuwenhuizen and D.Z. Freedman (North-Holland, Amsterdam, 1979); T. Yanagida, in *Proceedings of Workshop on Unified Theory and Baryon Number in the Universe*, edited by O. Sawada and A. Sugamoto (KEK, Tsukuba, 1979); R.N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 912 (1980).

[7] M.A. Luty, Phys. Rev. D **45**, 455 (1992).

[8] M. Plümacher, Z. Phys. C **74**, 549 (1997); Nucl. Phys. **B530**, 207 (1998); W. Büchmuller and M. Plümacher, Phys. Rep. **320**, 329 (1999); Int. J. Mod. Phys. A **15**, 5047 (2000).

[9] W. Buchmuller, P. Di Bari, and M. Plumacher, Nucl. Phys. **B643**, 367 (2002).

[10] J. Ellis, J.E. Kim, and D.V. Nanopoulos, Phys. Lett. **145B**, 181 (1984); J. Ellis, D.V. Nanopoulos, K.A. Olive, and S.-J. Rey, Astropart. Phys. **4**, 371 (1996).

[11] S. Sarkar, Rep. Prog. Phys. **59**, 1493 (1996).

[12] A.L. Maroto and A. Mazumdar, Phys. Rev. Lett. **84**, 1655 (2000); R. Kallosh, L. Kofman, A.D. Linde, and A. Von Proeyen, Phys. Rev. D **61**, 103503 (2000); G.F. Giudice, I.I. Tkachev, and A. Riotto, J. High Energy Phys. **08**, 009 (1999); **11**, 036 (1999); M. Bastero-Gil and A. Mazumdar, Phys. Rev. D **62**, 083510 (2000).

[13] R. Allahverdi, M. Bastero-Gil, and A. Mazumdar, Phys. Rev. D **64**, 023516 (2001).

[14] H.P. Nilles, M. Peloso, and L. Sorbo, J. High Energy Phys. **04**, 004 (2001).

[15] H.P. Nilles, K.A. Olive, and M. Peloso, Phys. Lett. B **522**, 304 (2001).

[16] R. Allahverdi, K. Enqvist, and A. Mazumdar, Phys. Rev. D **65**, 103519 (2002).

[17] G. Lazarides and Q. Shafi, Phys. Lett. B **258**, 305 (1991); T. Asaka, K. Hamaguchi, M. Kawasaki, and T. Yanagida, *ibid.* **464**, 12 (1999); Phys. Rev. D **61**, 083512 (2000).

[18] G.F. Giudice, M. Peloso, A. Riotto, and I.I. Tkachev, J. High Energy Phys. **08**, 014 (1999).
[19] K. Enqvist, S. Kasuya, and A. Mazumdar, Phys. Rev. Lett. **89**, 091301 (2002); Phys. Rev. D **66**, 043505 (2002).
[20] B.A. Campbell, S. Davidson, and K.A. Olive, Nucl. Phys. **B399**, 111 (1993).
[21] L. Boubekeur, S. Davidson, M. Peloso, and L. Sorbo, Phys. Rev. D **67**, 043515 (2003).
[22] H. Murayama and T. Yanagida, Phys. Lett. B **322**, 349 (1994); K. Hamaguchi, H. Murayama, and T. Yanagida, Phys. Rev. D **65**, 043512 (2002).
[23] Z. Berezhiani, A. Mazumdar, and A.P. Lorenzana, Phys. Lett. B **518**, 282 (2001).
[24] L. Bento and Z. Berezhiani, Phys. Rev. Lett. **87**, 231304 (2001).
[25] R. Allahverdi and A. Mazumdar, Phys. Rev. D **67**, 023509 (2003).
[26] E. Ma, M. Raidal, and U. Sarkar, Phys. Rev. D **60**, 076005 (1999).
[27] See, for example, T. Hambye, Nucl. Phys. **B633**, 171 (2002).
[28] M. Flanz, E.A. Paschos, and U. Sarkar, Phys. Lett. B **345**, 248 (1995); L. Covi, E. Roulet, and F. Vissani, *ibid.* **384**, 169 (1996); A. Pilaftsis, Phys. Rev. D **56**, 5431 (1997); W. Büchmuller and M. Plümacher, Phys. Lett. B **431**, 354 (1998).
[29] A. Pilaftsis, Int. J. Mod. Phys. A **14**, 1811 (1999).
[30] S. Davidson and S. Sarkar, J. High Energy Phys. **11**, 012 (2000); R. Allahverdi and M. Drees, Phys. Rev. D **66**, 063513 (2002); P. Jaikumar and A. Mazumdar, hep-ph/0212265.
[31] M. Fukugita and T. Yanagida, Phys. Rev. D **42**, 1285 (1990).
[32] K.S. Babu, B. Dutta, and R.N. Mohapatra, Phys. Rev. D **67**, 076006 (2003).