

# Sneutrino Leptogenesis at the Electroweak Scale

John Ellis

*TH Division, PH Department, CERN, 1211 Geneva 23, Switzerland*

Sin Kyu Kang

*School of Physics, Seoul National University, Seoul 151-734, Korea*

## Abstract

We propose an alternative mechanism for leptogenesis at the electroweak scale, through the decays of a left-handed sneutrino. This scenario may be realized in supersymmetric models with non-zero Majorana masses for the neutrino superfield that lead to mixing and mass splitting between the left-handed sneutrino and the corresponding antisneutrino. Soft supersymmetry breaking provides new sources of CP violation in sneutrino-antisneutrino mixing that can generate a lepton asymmetry in decays of the left-handed sneutrino. We show how the three Sakharov conditions for generating the observed baryon asymmetry of the Universe can be fulfilled in this restrictive framework.

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## I. INTRODUCTION

The baryon asymmetry of the universe may be expressed in terms of the ratio of the baryon density  $n_B$  to the entropy density  $s$  of the Universe. The agreement between astrophysical observations and primordial nucleosynthesis calculations constrains this ratio to the range at 95% C.L. [1]:

$$\frac{n_B}{s} \simeq (3.4 - 6.9) \cdot 10^{-10}, \quad (1)$$

which is in good agreement with the inferred from Cosmic Microwave Background data, particularly that provided by WMAP [2]. Almost 40 years ago, Sakharov explained how this baryon asymmetry could be obtained via a microphysical mechanism incorporating interactions that violate baryon and/or lepton number, violate C and CP, and drop out of equilibrium in the early Universe [3].

Since such interactions are present in the electroweak theory and its extensions, the baryon asymmetry might have originated at any energy above the electroweak scale. One of the most attractive and simple ways to realize the Sakharov mechanism is high-scale leptogenesis through the out-of-equilibrium decays of a heavy Majorana neutrino, since it arises naturally within the seesaw mechanism for generating light neutrino masses [4]. However, high-scale leptogenesis and the seesaw mechanism are difficult to test directly, since the natural scale of the interactions of the heavy Majorana neutrino is in the range  $10^{10}$  GeV to  $10^{15}$  GeV. This is impossible to reach by accelerators, except indirectly via renormalization effects on light particle masses, for example.

Baryogenesis at the electroweak scale itself is certainly a most attractive option[5], because it may be tested directly, if it can be achieved economically. However, this option is now excluded in the Standard Model, because the unsuccessful LEP searches require the Higgs boson to be too heavy for the electroweak phase transition to have been first-order, and the effective amount of CP violation is in any case very small [6]. On the other hand, an electroweak mechanism may be possible in the minimal supersymmetric extension of the Standard Model (MSSM), but at the price of a certain number of rather restrictive conditions [7]. It is therefore worthwhile to look for other scenarios for electroweak baryogenesis.

An additional problem with any supersymmetric model that postulates high temperatures in the early Universe is that of the gravitino abundance [8]. Thermal production of gravitinos at high temperatures in the very early Universe might have exceeded the bounds imposed

by primordial nucleosynthesis and the Cosmic Microwave Background [9]. Therefore, in the context of supersymmetry, it is doubly interesting to look for alternative baryogenesis scenarios that operate at lower energy scales below about a TeV, which might be tested directly in the foreseeable future.

It has recently been proposed that soft supersymmetry-breaking terms comprising bilinear and trilinear scalar couplings involving the right-handed sneutrino fields could be responsible for leptogenesis, an option referred to as soft leptogenesis [10]. In contrast to standard leptogenesis from heavy Majorana (s)neutrino decay, in soft leptogenesis the CP violation needed can be provided within the framework of a single sneutrino generation. However, even if this mechanism of leptogenesis could evade the gravitino problem, it is still difficult to probe in laboratory experiments, because of the high mass scale of the right-handed sneutrino.

In this paper, we propose an alternative leptogenesis mechanism via *decays of a left-handed sneutrino at a low energy scale*. This scenario may be realized in supersymmetric models with non-zero Majorana neutrino masses that lead to mixing and mass splitting between the left-handed sneutrino  $\tilde{\nu}$  and the corresponding antisneutrino  $\tilde{\bar{\nu}}$ . As was shown in [11, 12], the neutrino mass and the sneutrino mass splitting may be due to related  $\Delta L = 2$  interactions, linked by supersymmetry breaking. As we show below, the  $\tilde{\nu} - \tilde{\bar{\nu}}$  system may exhibit oscillatory behaviour analogous to the  $B - \bar{B}$  system. In addition, soft supersymmetry breaking sectors provide new sources of CP violation in  $\tilde{\nu} - \tilde{\bar{\nu}}$  mixing which can generate a lepton asymmetry in the sneutrino decays. We explore in this paper how the three Sakharov conditions [3] for generating the observed baryon asymmetry of the Universe can be fulfilled in this restrictive framework.

In the sense that supersymmetry breaking is the source of leptogenesis, our scenario is similar to soft leptogenesis, but differs in that it operates at the electroweak scale. If the scenario proposed does become a candidate for leptogenesis, it could in principle be tested in collider experiments via a sneutrino oscillation signal, though this may be difficult in practice. The lepton number is tagged in the decay of sneutrino by identifying the charge of the outgoing lepton, and a same-sign dilepton signal may be observable when the sneutrino-antisneutrino pairs decay into charged leptons. However, measurement of a lepton asymmetry would probably need more events than are likely to be provided at the LHC or ILC.

The outline of this paper is as follows. In the next section, we review how to generate the sneutrino mass splitting in a supersymmetric model with a heavy right-handed Majorana neutrino superfield and discuss how small the mass splitting can be. In section III, we examine how leptogenesis via the decay of the left-handed sneutrino can be realized and in section IV the CP asymmetry generated from the decay of the light sneutrino is considered. In section V, we discuss possible wash-out processes, which may be harmless for our leptogenesis scenario, and present our conclusions.

## II. LIGHT SNEUTRINO MASSES

We first review in more detail how to generate the sneutrino mass splitting in a supersymmetric model with non-zero Majorana neutrino masses, which could arise from a right-handed neutrino superfield. The relevant superpotential is

$$W = Y_\nu \widehat{H}_2 \widehat{L} \widehat{N} - \mu \widehat{H}_1 \widehat{H}_2 + \frac{1}{2} M \widehat{N} \widehat{N}, \quad (2)$$

where  $\widehat{L}, \widehat{N}, \widehat{H}_i (i = 1, 2)$  are superfields for lepton doublets, heavy right-handed neutrinos and Higgs fields, respectively. The  $D$  terms are the same as in the MSSM. The relevant terms in the soft supersymmetry-breaking scalar potential are given by

$$V_{soft} = m_L^2 \widetilde{\nu}^* \widetilde{\nu} + m_N^2 \widetilde{N}^* \widetilde{N} + (Y_\nu A_N H_2 \widetilde{\nu} \widetilde{N}^* + M B_N \widetilde{N} \widetilde{N} + H.c.), \quad (3)$$

where  $\widetilde{\nu}$  and  $\widetilde{N}$  are the light and heavy sneutrinos, respectively. From now on, we consider a single generation of  $\widehat{L}$  and  $\widehat{N}$ . After the electroweak symmetry is spontaneously broken, the light neutrino acquires a mass via the seesaw mechanism:  $m_\nu \simeq m_D^2/M$ , where  $m_D = Y_\nu v_2$  with  $v_2/\sqrt{2} = \langle H_2^0 \rangle$ . Defining

$$\begin{aligned} \widetilde{\nu}_1 &= (e^{i\phi/2} \widetilde{\nu} + e^{-i\phi/2} \widetilde{\nu}^*)/\sqrt{2}, & \widetilde{\nu}_2 &= -i(e^{i\phi/2} \widetilde{\nu} - e^{-i\phi/2} \widetilde{\nu}^*)/\sqrt{2}, \\ \widetilde{N}_1 &= (e^{i\phi'/2} \widetilde{N} + e^{-i\phi'/2} \widetilde{N}^*)/\sqrt{2}, & \widetilde{N}_2 &= -i(e^{i\phi'/2} \widetilde{N} - e^{-i\phi'/2} \widetilde{N}^*)/\sqrt{2}, \end{aligned} \quad (4)$$

one can separate the sneutrino mass-squared matrix into two blocks and then, to first order in  $1/M$ , the two light sneutrino eigenstates are  $\widetilde{\nu}_1$  and  $\widetilde{\nu}_2$ , with the following masses-squared,

$$m_{\widetilde{\nu}_{1,2}}^2 = m_L^2 + \frac{1}{2} m_Z^2 \cos 2\beta \mp \frac{1}{2} \Delta m_\nu^2, \quad (5)$$

where the mass-squared difference  $\Delta m_\nu^2 = m_{\widetilde{\nu}_2}^2 - m_{\widetilde{\nu}_1}^2$  is of order  $1/M$ . Here, we can remove the phases in the superpotential (2) by rotating superfields and the relative phase between

$A_N$  and  $B_N$  with a R-rotation, but there remains an unremovable phase, which makes  $\Delta m_{\tilde{\nu}}^2$  complex. The sneutrino mass splitting is then easily computed by using the relation  $\Delta m_{\tilde{\nu}}^2 = 2m_{\tilde{\nu}}\Delta m_{\tilde{\nu}}$  [11], where  $m_{\tilde{\nu}} = \frac{1}{2}(m_{\tilde{\nu}_1} + m_{\tilde{\nu}_2})$  is the average of the light sneutrino masses and

$$\Delta m_{\tilde{\nu}} \simeq \frac{2m_{\nu}(A_N - \mu \cot \beta - B_N)}{m_{\tilde{\nu}}}. \quad (6)$$

We note that  $\mu$ ,  $A_N$  and  $m_{\tilde{L}}$  are of the same order as the electroweak scale, whereas  $M$ ,  $m_{\tilde{N}}$  and  $B_N$  are soft-supersymmetry breaking parameters associated with the  $SU(2) \times U(1)$  singlet superfield  $\widehat{N}$ , and may be much larger than the electroweak scale. As one can see from (6), if  $B_N \gg m_Z$ , the sneutrino mass splitting is significantly enhanced, whereas it is of the same order as the neutrino mass when  $B_N \sim \mathcal{O}(m_Z)$ . This mass splitting may in principle be probed through the sneutrino-antisneutrino oscillation which would result in a same-sign dilepton signal. However, to have an observable rate for the same-sign dilepton signal, the ratio of the mass splitting to the sneutrino decay width should be large, namely,  $\Delta m_{\tilde{\nu}}/\Gamma_{\tilde{\nu}} \geq 1$ , and the sneutrino branching ratio into a charged lepton should also be large.

### III. CONDITIONS FOR LEPTOGENESIS

We now examine how leptogenesis via the decay of the light sneutrino can be realized in this framework. As is well known, for a successful mechanism of leptogenesis, we need lepton-number- and CP-violating interactions that should be out of equilibrium, so that the asymmetry generated is not automatically suppressed. In this scenario, as mentioned before, *sneutrino-antisneutrino mixing is a direct manifestation of lepton-number violation*, and the *soft supersymmetry-breaking terms may well provide a suitable new source of CP violation*.

We must then impose the *out-of-equilibrium condition for the decay of the light sneutrino*, which is given by

$$\Gamma_{\tilde{\nu}} < H \simeq 1.7\sqrt{g_*}\frac{T^2}{M_P}, \quad \text{at } T = m_{\tilde{\nu}}, \quad (7)$$

where  $\Gamma_{\tilde{\nu}}$  is the decay rate of the sneutrino,  $H$  is the Hubble constant at the decay epoch, and  $M_P$  is the Planck scale. The parameter  $g_*$  is the effective number of massless degrees of freedom, which takes the value  $g_* = 225$  for the MSSM with one generation of right-handed neutrinos.

As well as the out-of-equilibrium condition (7), we must require that the electroweak *sphaleron interactions are still in thermal equilibrium* at the time the lepton asymmetry is generated, so that they can convert the lepton asymmetry partially into a baryon asymmetry of the Universe. The temperature at which the sphaleron interactions freeze out depends on how the electroweak phase transition occurs. It is known that the sphaleron interactions freeze out at the critical temperature of electroweak phase transition if it is strongly first-order, whereas the freeze-out temperature may become lower than the critical temperature if the transition is second-order or weakly first-order. In the latter case, the sphaleron interactions freeze out at the temperature at which the sphaleron transition rate  $\Gamma_{sph}$  becomes equal to the expansion rate of the Universe[13]. It has been found in numerical simulations that the sphaleron interactions are effective as long as  $T \geq 200$  GeV. Therefore, for a successful mechanism of leptogenesis, we require

$$5 \times 10^{-4} \text{ eV} \lesssim \Gamma_{\tilde{\nu}} \lesssim 1.3 \times 10^{-4} \left( \frac{m_{\tilde{\nu}}(\text{GeV})}{100\text{GeV}} \right)^2 \text{ eV}, \quad (8)$$

where the first condition follows from the sphaleron equilibrium condition and the second from (7) for  $T = m_{\tilde{\nu}}$ . As seen in (8), a scenario for leptogenesis can be successful only when  $m_{\tilde{\nu}} \gtrsim 200$  GeV.

#### IV. CP-VIOLATING SNEUTRINO DECAYS AND MIXING

In order to discuss the decays, mixing and CP-violating asymmetries for light sneutrinos, we first integrate out the heavy right-handed neutrino superfield, following which the superpotential for the light sneutrino is:

$$W = \lambda_{ij} \widehat{\nu}_i \widehat{l}_R^c \widehat{H}_1 + \frac{\kappa_{ij}}{M} \widehat{H}_2 \widehat{\nu}_i \widehat{\nu}_j \widehat{H}_2^T, \quad (9)$$

where  $\widehat{\nu}$ ,  $\widehat{l}_R^c$ ,  $\widehat{H}_1$  denote the light neutrino, the right-handed charged lepton and the charged Higgs superfields, respectively, and  $\lambda_{ij}$  stands for the Yukawa couplings of the lepton sector, which are given by  $\lambda_{ij} = -gm_{i_{ij}}/\sqrt{2}M_W \cos \beta$ . The soft supersymmetry-breaking terms involving the light sneutrinos  $\tilde{\nu}$  are;

$$-L_{soft} = m_{\tilde{L}}^2 \tilde{\nu}^* \tilde{\nu} + \Delta m_{\tilde{\nu}}^2 \tilde{\nu} \tilde{\nu} + \lambda A_{\nu} \tilde{\nu} \tilde{l}_R H_1. \quad (10)$$

We note that the effects of the right-handed neutrino superfields are absorbed into the seesaw term in (9) and the sneutrino mass splitting  $\Delta m_{\tilde{\nu}}^2$  discussed previously. The sneutrino

interaction Lagrangian is then given by

$$L = \tilde{\nu}(\lambda\widetilde{H}_1 l_R + m_\nu \lambda^* \tilde{l}_{Ri}^* H_1^* + A_\nu \lambda \tilde{l}_R H_1 + g' Z_{iZ} \tilde{\chi}_i^0 \nu + g V_{11} \tilde{\chi}^+ l_L^-) + h.c., \quad (11)$$

where we have considered a single sneutrino generation and  $m_\nu = \kappa v_2^2/M$  after the electroweak symmetry is broken.

It should be noted that the Higgs field  $H_1$  is decomposed into the physical Higgs sector and the Goldstone boson sector. Selecting the physical Higgs sector in  $H_1$  is equivalent to replacing  $H_1$  by  $H^- \sin \beta$ . Its Yukawa coupling therefore becomes  $-gm_l \tan \beta / \sqrt{2} M_W$ , where  $m_l$  is the charged lepton mass. In the Lagrangian for a single generation of sneutrino, there is a physical CP-violating phase. With superfield rotations and an R-rotation we can eliminate CP phases in the parameters  $\Delta m_{\tilde{\nu}}^2$ ,  $\lambda$  and  $m_\nu$ , but then the CP-violating phase in  $A_\nu$  cannot be removed.

As observed above in (8), the mechanism of low-energy leptogenesis proposed here could be successful only if  $\Gamma_{\tilde{\nu}}$  is of order  $10^{-4}$  eV, so it is important to discuss how such a low decay rate might be achieved. If  $m_{\tilde{\chi}_1^0} < m_{\tilde{\chi}^+}, m_{\tilde{\nu}}$ , which is the ordering of masses generally expected in the MSSM with universal soft supersymmetry-breaking scalar masses (the CMSSM), the dominant sneutrino decays are those into two-body final states. As shown in [11], the typical size of  $\Gamma_{\tilde{\nu}}$  for such two-body decays is  $\mathcal{O}(\text{GeV})$ , which is too large to achieve successful leptogenesis.

On the other hand, if  $m_{\tilde{\nu}} < m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}^+}$  and no two-body sneutrino decay channels are open, three-body sneutrino decays will dominate. This mass ordering is possible if the gravitino is the lightest supersymmetric particle, and is also allowed by cosmology and the standard accelerator constraints [14]. The following chargino- and neutralino-mediated three-body decays are then generally dominant:  $\tilde{\nu}_l \rightarrow l^- \tilde{\tau}_R^+ \nu_\tau$  and  $\tilde{\nu}_l \rightarrow \nu_l \tilde{\tau}^\pm \tau^\mp$ , assuming that decays into final states containing lighter sneutrinos can be neglected. Assuming that the lightest neutralino is dominated by its bino component, the rates for the chargino- and neutralino-mediated sneutrino decays (see Fig. 1(a) for an illustration) are given by [11]

$$\Gamma(\tilde{\nu}_l \rightarrow l^- \tilde{\tau}_R^+ \nu_\tau) = \frac{g^4 m_{\tilde{\nu}}^3 m_\tau^2 \tan^2 \beta f_c(m_\tau^2/m_{\tilde{\nu}}^2)}{3 \times 2^9 \pi^3 (m_W^2 \sin 2\beta - M_2 \mu)^2}, \quad (12)$$

$$\Gamma(\tilde{\nu}_l \rightarrow \nu_l \tilde{\tau}_R^\pm \tau^\mp) = \frac{g'^4 m_{\tilde{\nu}}^5 f_n(m_{\tilde{\tau}_R}^2/m_{\tilde{\nu}}^2)}{3 \times 2^{10} \pi^3 M_1^4}, \quad (13)$$

where the  $M_i$  are gaugino mass parameters,  $f_c(x) \equiv (1-x)(1+10x+x^2) + 3x(1+x) \log x^2$  and  $f_n(x) \equiv 1 - 8x + 8x^3 - x^4 - 6x^2 \log x^2$ .

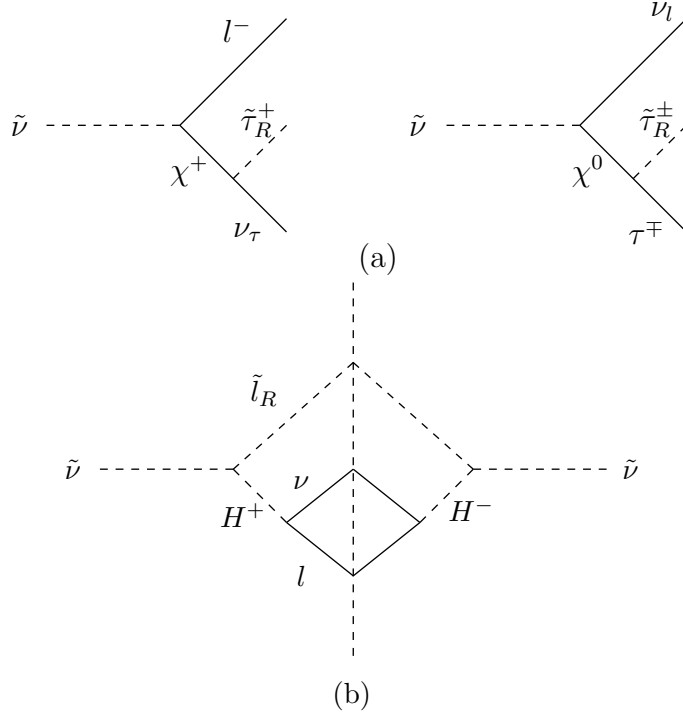


FIG. 1: (a) Diagrams for three-body sneutrino decay via charginos and neutralinos, and (b) a diagram contributing to  $\Gamma_{12}$ .

In order to obtain  $\Gamma \simeq 5 \times 10^{-4}$  eV, one must require that the  $\tilde{\nu}$  and  $\tilde{\tau}_R$  be nearly degenerate. As an example, for  $m_{\tilde{\tau}_R}^2/m_{\tilde{\nu}}^2 = 0.95$ ,  $m_{\tilde{\nu}} \simeq 500$  GeV, and  $M_1 \simeq 350$  GeV, we indeed find  $\Gamma \simeq 5 \times 10^{-4}$  eV. This near-degeneracy between the  $\tilde{\tau}_R$  and  $\tilde{\nu}$  is not possible in the CMSSM with universal soft scalar masses  $m_0$  and gaugino masses  $m_{1/2}$  at the GUT scale, in which the mass difference between the  $\tilde{\tau}_R$  and  $\tilde{\nu}$  is approximately  $m_{\tilde{\nu}}^2 - m_{\tilde{\tau}_R}^2 \simeq 0.4m_{1/2}^2 + 0.27 \cos 2\beta m_Z^2$ . According to the constraints on the parameter space  $(m_0, m_{1/2})$  from the cosmological observation,  $b \rightarrow s\gamma$  and muon  $g - 2$ , the favoured region of  $m_{1/2}$  is somewhat large, namely  $m_{1/2} \geq 300$  GeV, and the  $\tilde{\tau}_R$  is considerably lighter than the  $\tilde{\nu}$  [14] in the CMSSM. Thus, we need a non-universal boundary condition on the soft scalar masses at the GUT scale. There is no known theoretical or phenomenological reason why the soft supersymmetry-breaking  $\tilde{\tau}_R$  and  $\tilde{\nu}$  masses should not differ at the GUT scale. At the electroweak scale,  $m_{\tilde{\tau}_R}$  and  $m_{\tilde{\nu}}$  can be written as

$$m_{\tilde{\tau}_R} \simeq m_{0R}^2 + 0.15m_{1/2}^2 + \sin^2 \theta_W \cos 2\beta m_Z^2, \quad (14)$$

$$m_{\tilde{\nu}} = m_{0L}^2 + 0.54m_{1/2}^2 + 1/2 \cos 2\beta m_Z^2, \quad (15)$$

and we can make the physical  $m_{\tilde{\tau}_R}$  and  $m_{\tilde{\nu}}$  nearly degenerate by choosing the soft scalar



masses  $m_{0R}$  and  $m_{0L}$  appropriately.

We now consider mass mixing and CP violation in the  $\tilde{\nu} - \tilde{\nu}^*$  system. Since it is analogous to the  $K^0 - \bar{K}^0$  and  $B^0 - \bar{B}^0$  systems, we can calculate mixing and CP asymmetries using the formulae derived previously for these meson systems [15]. The evolution of the system is determined by a Hamiltonian  $H = \widehat{M} - i\widehat{\Gamma}/2$  where, to leading order in the soft terms,

$$\widehat{M} = m_{\tilde{\nu}} \begin{pmatrix} 1 & \frac{\Delta m_{\tilde{\nu}}}{m_{\tilde{\nu}}} \\ \frac{\Delta m_{\tilde{\nu}}}{m_{\tilde{\nu}}} & 1 \end{pmatrix}, \quad (16)$$

$$\widehat{\Gamma} = \begin{pmatrix} \Gamma & \Gamma_{12} \\ \Gamma_{12}^* & \Gamma \end{pmatrix} \quad (17)$$

where  $\Gamma$  is the total  $\tilde{\nu}$  decay width. The dominant part of the off-diagonal component  $\Gamma_{12}$  can be obtained by considering the imaginary part of the loop diagram shown in Fig. 1(b). The analytic expression for  $\Gamma_{12}$  is given by

$$\Gamma_{12} = \frac{g^4 \tan^4 \beta m_l^4 m_\nu A_\nu}{512\pi^3 M_W^4 m_{\tilde{\nu}}^3} I \quad (18)$$

where

$$I = \int_0^{(m_{\tilde{\nu}} - m_l)^2} ds \frac{s \sqrt{s^2 - 2(m_{\tilde{\nu}}^2 + m_l^2)s + (m_{\tilde{\nu}}^2 - m_l^2)^2}}{(s - m_{H^-}^2)}. \quad (19)$$

The eigenvectors of the Hamiltonian  $H$  are then given by

$$\tilde{\nu}_L = p\tilde{\nu} + q\tilde{\nu}^*, \quad \tilde{\nu}_H = p\tilde{\nu} - q\tilde{\nu}^*, \quad (20)$$

where the parameters  $p$  and  $q$  are related by

$$\left(\frac{q}{p}\right)^2 = \frac{\widehat{M}_{12}^* - \frac{i}{2}\widehat{\Gamma}_{12}^*}{\widehat{M}_{12} - \frac{i}{2}\widehat{\Gamma}_{12}}. \quad (21)$$

It is appropriate for cosmology to consider an initial state at  $t = 0$  with equal densities of  $\tilde{\nu}$  and  $\tilde{\nu}^*$ . At a later time  $t$ , the state will have evolved into

$$\tilde{\nu}(t) = g_+(t)\tilde{\nu}(0) + \frac{q}{p}g_-(t)\tilde{\nu}^*(0), \quad \tilde{\nu}^*(t) = \frac{p}{q}g_-(t)\tilde{\nu}(0) + g_+(t)\tilde{\nu}^*(0), \quad (22)$$

$$g_+(t) = e^{-im_{\tilde{\nu}}t} e^{-\Gamma t/2} \cos(\Delta m_{\tilde{\nu}} t/2), \quad g_-(t) = e^{-im_{\tilde{\nu}}t} e^{-\Gamma t/2} \sin(\Delta m_{\tilde{\nu}} t/2). \quad (23)$$

Here  $\Delta m_{\tilde{\nu}} \equiv m_{\tilde{\nu}_2} - m_{\tilde{\nu}_1}$  and we have neglected  $\Delta\Gamma$  with respect to  $\Delta m_{\tilde{\nu}}$ . We can then

compute the total integrated lepton asymmetry, defined by

$$\varepsilon = \frac{\sum_f \int_0^\infty dt [\Gamma(\tilde{\nu}(t) \rightarrow f) + \Gamma(\tilde{\nu}^*(t) \rightarrow f) - \Gamma(\tilde{\nu}(t) \rightarrow \bar{f}) - \Gamma(\tilde{\nu}^*(t) \rightarrow \bar{f})]}{\sum_f \int_0^\infty dt [\Gamma(\tilde{\nu}(t) \rightarrow f) + \Gamma(\tilde{\nu}^*(t) \rightarrow f) + \Gamma(\tilde{\nu}(t) \rightarrow \bar{f}) + \Gamma(\tilde{\nu}^*(t) \rightarrow \bar{f})]} \quad (24)$$

$$= \frac{1}{2} \left( \left| \frac{q}{p} \right|^2 - \left| \frac{p}{q} \right|^2 \right) \frac{\int_0^\infty dt |g_-|^2}{\int_0^\infty dt (|g_+|^2 + |g_-|^2)}, \quad (25)$$

where  $f$  stands for a final state with lepton number equal to one and  $\bar{f}$  is its conjugate.

Calculating (21) in the limit  $\widehat{\Gamma}_{12} \ll \widehat{M}_{12}$ , we find

$$\begin{aligned} \left| \frac{q}{p} \right|^2 &\simeq 1 - \text{Im} \left( \frac{\widehat{\Gamma}_{12}}{\widehat{M}_{12}} \right) \\ &\simeq 1 - \text{Im} \left( \frac{g^4 \tan^4 \beta m_l^4 m_\nu A_\nu I}{512 \pi^3 M_W^4 m_\nu^3 \Delta m_{\tilde{\nu}}} \right). \end{aligned} \quad (26)$$

We notice that  $A_\nu$  is the only complex parameter in this expression. Performing the time integral, we find

$$\frac{\int_0^\infty dt |g_-|^2}{\int_0^\infty dt (|g_+|^2 + |g_-|^2)} = \frac{(\Delta m_{\tilde{\nu}})^2}{2(\Gamma^2 + (\Delta m_{\tilde{\nu}})^2)}. \quad (27)$$

Thus, we obtain the following final expression for the CP asymmetry:

$$\varepsilon = - \frac{(\Delta m_{\tilde{\nu}})^2}{2(\Gamma^2 + (\Delta m_{\tilde{\nu}})^2)} \frac{g^4 \tan^4 \beta m_l^4 m_\nu \text{Im} A_\nu I}{512 \pi^3 M_W^4 m_\nu^3 \Delta m_{\tilde{\nu}}}. \quad (28)$$

The baryon asymmetry is then given by [16]

$$\frac{n_B}{s} = - \left( \frac{24 + 4n_H}{66 + 13n_H} \right) \varepsilon \eta Y_{\tilde{\nu}}^{eq}. \quad (29)$$

The first factor takes into account the reprocessing of the  $B - L$  asymmetry by sphaleron transitions, with the number of Higgs doublets  $n_H$  equal to 2, and  $Y_{\tilde{\nu}}^{eq} = 45\zeta(3)/(\pi^4 g_*)$  is the equilibrium sneutrino density in units of the entropy density for temperatures much larger than  $m_{\tilde{\nu}}$ . Therefore, we obtain

$$\frac{n_B}{s} = -8.6 \times 10^{-4} \varepsilon \eta. \quad (30)$$

The efficiency factor  $\eta$  includes effects caused by the sneutrino density being smaller than the equilibrium density and wash-out due to the decays not being completely out of equilibrium. The value of  $\eta$  must be obtained by integrating the relevant Boltzmann equations, but may plausibly be of order unity.

For the value of  $n_B/s$  given in (1), we see that that our mechanism should yield  $|\varepsilon\eta| \simeq 10^{-6} - 10^{-7}$ . Let us estimate how we can achieve this desirable amount of  $|\varepsilon|$ . We infer from recent experimental results that a neutrino mass of order  $10^{-2}$  eV is possible. Then, we see from (5) that  $\Delta m_{\tilde{\nu}} \simeq 10^{-2}$  eV in the case that the parameters  $A_N, B_N, \mu$  are of the same order as  $m_{\tilde{\nu}}$ . In this case, the size of the quantity in (27) is of order one. We can also take the parameter  $\text{Im}A_{\nu}$  to be the same order of  $m_{\tilde{\nu}}$ . Moreover, we observe that the value of  $\Gamma_{12}$  can be enhanced by taking large  $\tan\beta$ . Numerically, we find a magnitude of  $|\varepsilon|$  in (28) of order  $10^{-6} - 10^{-7}$  when we take  $m_{\tilde{\nu}} \simeq 500\text{GeV}$ ,  $m_H \simeq 1\text{TeV}$ ,  $m_l = m_{\tau}$  and  $\tan\beta \simeq 30$ , and a larger value of  $|\varepsilon|$  could be obtained if needed.

Finally, we note that the large sneutrino-antisneutrino mixing which is expected because  $\Delta m_{\tilde{\nu}} \gg \Gamma_{\tilde{\nu}}$  might lead to an observable like-sign dilepton signal. This could provide an characteristic collider signature of our scenario. On the other hand, the amount of CP violation is likely to be too small to be observable in the near future.

## V. DISCUSSION

The scenario proposed here needs further investigation. In particular, we recall that successful leptogenesis requires a departure from thermal equilibrium for wash-out processes, as well lepton-number-violating decays. The dominant wash-out processes are scattering reactions mediated by sneutrinos or charginos and neutralinos. Since our scenario calls for relatively heavy charginos and neutralinos, the scattering rates mediated by charginos and neutralinos are suppressed. Additionally, scattering reactions mediated by sneutrinos can occur via the vertices involving charginos or neutralinos. However, the rates of these scattering processes are suppressed due to the Boltzmann suppression of the number densities of the charginos and neutralinos. A more accurate estimation of the lepton asymmetry that survives these wash-out processes could be obtained by solving directly the full Boltzmann equations for the system with different sparticle masses, as has been done in many detailed works considering high-scale leptogenesis [16]. A similarly detailed study goes beyond the scope of this paper, and will be undertaken elsewhere.

In conclusion, an alternative of leptogenesis through decays of a left-handed sneutrino at a low energy scale has been proposed. This scenario may be realized in supersymmetric models with non-zero Majorana masses of neutrino superfield that lead to mixing and mass splitting

between the left-handed sneutrino and the corresponding antisneutrino. Soft supersymmetry breaking sectors provide new sources of CP violation in sneutrino-antisneutrino mixing which can generate a lepton asymmetry in decays of the left-handed sneutrino. We have shown that the three Sakharov conditions [3] for the observed baryon asymmetry of the Universe can be fulfilled in this restrictive framework.

## VI. ACKNOWLEDGEMENT

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