Received: *June 14, 2000*, Accepted: *July 6, 2000*

The minimal supersymmetric leptogenesis*

Takeo Moroi

School of Natural Sciences, Institute for Advanced Study
Princeton, New Jersey 08540
E-mail: moroi@ias.edu

Hitoshi Murayama

Theoretical Physics Group
Ernest Orlando Lawrence Berkeley National Laboratory
University of California, Berkeley, California 94720, and
Department of Physics, University of California
Berkeley, California 94720
E-mail: murayama@thsrv.lbl.gov

Abstract: We reanalyze the supersymmetric leptogenesis model by one of the authors (H.M.) and Yanagida based on $\tilde{L} = H_u$ flat direction in detail. We point out that the appropriate amount of baryon asymmetry can be generated in this model with the neutrino mass matrix consistent with the atmospheric neutrino oscillation and solutions to the solar neutrino problem, preferably the small angle MSW solution. The reheating temperature can be low enough to avoid the cosmological gravitino problem. This model is the minimal one because it does not rely on any new physics beyond supersymmetry and Majorana neutrino masses.

Keywords: Supersymmetric Effective Theories, Baryogenesis, Neutrino Physics.

*This work was supported in part by the U.S. Department of Energy under Contracts DE-AC03-76SF00098, in part by the National Science Foundation under grants PHY-95-13835 and PHY-95-14797. TM was also supported by the Marvin L. Goldberger Membership. H.M. was also supported by Alfred P. Sloan Foundation.

Contents

1. Introduction

Cosmic baryon asymmetry has been one of the biggest puzzles in modern cosmology. It is hoped to understand it in terms of microscopic physics and evolution of Universe, if the well-known Sakharov's three conditions are satisfied: (1) existence of baryon-number violation, (2) existence of CP violation, and (3) departure from thermal equilibrium. A stringent constraint on the baryogenesis models is posed by the apparent lack of baryon-number violation. Recent data from Soudan-II and Super-Kamiokande experiments [1] made the constraints on proton and neutron decay yet more stronger. Especially given the constraint on $p \to e^+\pi^0$, which requires the mass of the X boson in SU(5) grand unified theories to be above 10^{15} GeV or so, it is becoming increasingly difficult to construct models of baryogenesis with explicit baryon-number violation. On the other hand, the recent evidence for oscillation in atmospheric neutrinos from Super-Kamiokande [2], Soudan-II [3] and MACRO [4] experiments strongly suggest small but finite neutrino masses with $\Delta m^2 = 10^{-3}$–$10^{-2}\,\mathrm{eV}^2$. Such small mass scales are unlikely to arise for Dirac neutrinos,[1] while are naturally understood in terms of so-called seesaw mechanism [5] which in turn requires Majorana neutrinos. Therefore, given the theoretical prejudice, there is a strong evidence for lepton-number violation at the mass scale of the right-handed neutrinos $M_R \sim 10^{14}$–10^{15} GeV.

[1] An alternative possibility was pointed out recently using large extra dimensions with right-handed neutrinos in the bulk [6].

Fukugita and Yanagida [7] suggested that the lepton-number violation may be enough for baryogenesis because the electroweak anomaly violates $B + L$ [8] and, once lepton number is violated, so is baryon number. If a finite lepton asymmetry is created in the Early Universe, it can be partially converted to baryon asymmetry without a need for an explicit baryon-number violating process which could lead to too-rapid proton decay. This scenario, called leptogenesis, is hence probably the most attractive possibility of baryogenesis in the current situations. The thermally produced right-handed neutrinos decay out of equilibrium and the CP-violation in the neutrino Yukawa matrix causes a difference between the decay rates into leptons and anti-leptons in this model. The lightest right-handed neutrino is typically as heavy as 10^{10} GeV in many models and therefore the reheating temperature after inflation needs to be higher than that (see [9] for a recent discussion).

On the other hand, the large hierarchy between the electroweak scale and the scale of lepton number violation M_R would require a stabilization mechanism for the hierarchy. Supersymmetry is the best solution to this problem currently available. Therefore it is natural and timely to ask what the minimal scheme is for the supersymmetric leptogenesis.

Supersymmetry, however, causes a cosmological problem when the reheating temperature after the primordial inflation is too high. The gravitinos can be produced thermally whose abundance is approximately proportional to the reheating temperature. For a typical mass range $m_{3/2} \sim 500$ GeV favored to stabilize the hierarchy,[2] the produced gravitinos decay after the Big-Bang Nucleosynthesis and they dissociate light elements. To keep the theory and observation consistent, there is an upper bound on the gravitino abundance, and hence on the reheating temperature of about $T_{RH} \lesssim 10^9$ GeV for this mass range (for the latest analysis, see [11]). This constraint makes it somewhat difficult to produce right-handed neutrino states as required in the original model [7].[3]

In this letter, we study the minimal supersymmetric leptogenesis. Our definition of minimal is the particle content of the Minimal Supersymmetric Standard Model together with suggested neutrino masses as effective dimension-five operators. This way, we avoid reliance on explicit grand-unified models or seesaw models. The only dependence on physics beyond the standard model is supersymmetry and neutrino masses. We show that leptogenesis is possible à la Affleck-Dine [14] within this minimal framework as originally suggested by one of the authors (H.M.) and Yanagida [15], within the current experimental situations. The reheating temperature can be low enough to avoid the gravitino problem because the model does not require the production of right-handed neutrinos.

[2] Anomaly mediation of supersymmetry breaking [10] may allow much higher gravitino mass keeping most superparticles light, and hence a higher reheating temperature.

[3] One can get around this problem, however, if the right-handed neutrinos are produced non-thermally, i.e. by the decay of the inflaton [12] or preheating [13].

2. The model

We start with the effective dimension-five operators in the superpotential

$$W = \frac{m_{\nu_i}}{2v_u^2}(L_i H_u)(L_i H_u) = \frac{1}{2M_i}(L_i H_u)(L_i H_u)\,, \tag{2.1}$$

which are necessary to incorporate finite neutrino masses (and hence lepton-number violation) to the Minimal Supersymmetric Standard Model. The vacuum expectation value for H_u is $v_u = \langle H_u \rangle = (174\,\mathrm{GeV})\,\sin\beta$ and we assume a moderately large $\tan\beta \gtrsim 5$ throughout the paper.[4] The index i runs over three mass eigenstates of neutrinos. It will turn out that the relevant direction is most likely the lightest neutrino as we will see later and henceforth suppress the index i. Along the flat direction $\tilde{L} = H_u = \phi/\sqrt{2}$, (the factor $\sqrt{2}$ is necessary to ensure the canonical kinetic term: $L^* L + H_u^* H_u = \phi^* \phi$) the potential is given by

$$V = m^2|\phi|^2 + \frac{A}{8M}(\phi^4 + \phi^{*4}) + \frac{1}{4M^2}|\phi|^6\,. \tag{2.2}$$

Here, the first two terms are from SUSY breaking and we assume $m \sim A \sim 100\,\mathrm{GeV}$–$1\,\mathrm{TeV}$, while the last term is from superpotential given in eq. (2.1). The mass scale M does not necessarily be the mass scale of right-handed neutrinos; it can actually be much higher. Our philosophy is not to discuss the precise origin of such operator, but rather use this effective operator only.

3. Initial amplitude

Since this direction is D-flat and also F-flat in the $m_\nu \to 0$ limit, it can have a large amplitude at the end of the primordial inflation. It turns out that the resulting baryon asymmetry does not depend on the precise value of the initial amplitude. All we need is that the initial amplitude is larger than $\sqrt{mM}$ as we will see later. We would like to discuss various mechanisms to generate a large amplitude proposed in the literature. We have the expansion rate of the Universe during the primordial inflation of $H_{\rm inf} \sim 10^{12}$–$10^{14}\,\mathrm{GeV}$ as suggested by popular models [16].

To discuss the amplitude in the Early Universe, one needs to pay special attention to the fact that the expanding Universe itself breaks global supersymmetry and may modify the scalar potential [18]. The explicit form of the supersymmetry breaking, however, depends on the details of the Kähler potential and hence is model dependent. Therefore we consider three possibilities in this letter. (1) No modification from the flat-space potential, which occurs at the tree-level of the no-scale type supergravity [19], (2) loop-level modifications from the flat-space potential [19], and

[4] For smaller $\tan\beta \gtrsim 1.4$ allowed by the perturbative top Yukawa coupling, the resulting lepton asymmetry changes only up to a factor of two.

(3) tree-level modifications from the flat-space potential [18]. In all cases, we find it is easy to have an initial amplitude necessary for the model: $\phi_{\rm inf} \gg \sqrt{mM}$ as we will see later.

In case (1), one idea to generate a large initial amplitude is to start with a "chaotic" initial condition [17],[5] where all scalar fields start with large values. During the course of inflation, most of the fields are diluted exponentially due to their potential, but some remain if their potential is flat enough. Along our flat direction, the potential is eventually dominated by $|\phi|^6$ term once the amplitude becomes small enough whatever the physics at high energy scale $\gg M$ is. The evolution is given by the slow-rolling approximation, and we can solve the equation of motion

$$3H_{\rm inf}\dot{\phi} + \frac{3}{4M^2}|\phi|^4\phi = 0\,. \tag{3.1}$$

For a more-or-less constant expansion rate $H = H_{\rm inf}$ during the inflation, its solution is

$$\frac{1}{|\phi|^4(t)} - \frac{1}{|\phi|^4(0)} = \frac{1}{M^2 H_{\rm inf}} t\,. \tag{3.2}$$

Therefore with an e-folding of $N = H_{\rm inf}t$,

$$|\phi|^4(t) = \frac{1}{|\phi|^{-4}(0) + N/(M^2 H_{\rm inf}^2)}\,. \tag{3.3}$$

Clearly, the amplitude at the end of inflation is almost predicted as long as $|\phi(0)|$ is large enough,

$$|\phi|^4(t) = \frac{M^2 H_{\rm inf}^2}{N}\,. \tag{3.4}$$

The e-folding N required to solve the horizon and flatness problems is $N \gtrsim 100$ [20].

The cases (2) and (3) use a possible negative mass squared of order $-H_{\rm inf}^2/16\pi^2$ [19] or $-H_{\rm inf}^2$ [18]. Minimizing

$$V = -cH_{\rm inf}^2|\phi|^2 + \frac{1}{4M^2}|\phi|^6\,, \tag{3.5}$$

with $c \sim O(1)$ or $O(1/16\pi^2)$, we find $-cH_{\rm inf}^2\phi + (3/4M^2)|\phi|^4\phi = 0$, or $|\phi|^4 = 4cH_{\rm inf}^2M^2/3$.

Both above arguments are classical. One may worry that the quantum fluctuation during the de Sitter expansion of Universe may modify the estimate. An effective Hawking-Gibbons temperature is given by $T_{HG} = H_{\rm inf}/2\pi$ and the fields acquire expectation values which correspond to the potential energy density of the order of the T_{HG}^4 [20]. This effect suggests $|\phi|^6 \sim M^2 H_{\rm inf}^4$, which is typically smaller than the other estimates for the range of M of our interest (see below). This smallness of the quantum fluctuation therefore justifies the classical treatment above for the flat direction because the fluctuation around the classical value can be neglected.

[5] We do not need $V \sim M_{\rm Planck}^4$ as suggested in [17], however.

Based on the above discussions, we take an order of magnitude estimate of the amplitude of ϕ at the end of the inflation $t_{\rm inf}$:

$$\phi_{\rm inf}^4 \equiv |\phi(t_{\rm inf})|^4 \sim M^2 H_{\rm inf}^2\,, \tag{3.6}$$

or slightly smaller for the rest of the analysis. The result turns out not to depend on the precise value of the initial amplitude; the only requirement is that $\phi_{\rm inf} \gg \sqrt{mM}$.

4. Evolution

Given the above initial condition, we follow the cosmological evolution after the end of the inflation. While the inflaton oscillates around the minimum of its potential, the flat direction ϕ gets diluted. For case (1), the equation of motion of the ϕ when the potential is dominated by ϕ^6 term is given by

$$\ddot{\phi} + \frac{2}{t}\dot{\phi} + \frac{3}{4}\frac{\phi^5}{M^2} = 0\,, \tag{4.1}$$

assuming the matter dominated expansion $H = 2/3t$. For case (2), there is an additional term of order $-\frac{1}{16\pi^2 t^2}\phi$ in the l.h.s. of the equation, and similarly $-\frac{1}{t^2}\phi$ for case (3). Despite the apparent difference between three cases, ϕ turns out to track $\sqrt{MH}$ until the expansion rate slows down to $H \sim m$ as we will explain below.

We first discuss case (1) with flat-space potential because of its simplicity. The behavior of $|\phi|$ depends on which term dominates its potential, since the potential changes its shape at

$$|\phi| \sim \sqrt{mM} \equiv \phi_{LG}\,. \tag{4.2}$$

At the time when $|\phi| \sim \phi_{LG}$, the lepton number gets fixed (we refer to this point as LG, for LeptoGenesis). For $|\phi| \gtrsim \phi_{LG}$, $|\phi|^6$ term dominates the potential and the equation of motion (4.1) should be used. Using the slow-roll approximation $|\ddot{\phi}| \ll |2\dot{\phi}/t|$, the solution is given by

$$\phi(t)^4 = \left[\phi_{\rm inf}^{-4} + 3\frac{t^2 - t_{\rm inf}^2}{4M^2}\right]^{-1}, \tag{4.3}$$

and it is easy to check that the motion is critically damped when $\phi_{\rm inf} \lesssim \sqrt{MH}$. Therefore, independent of the precise initial value of $\phi_{\rm inf}$, the motion starts only when $\phi_{\rm inf} \sim \sqrt{MH}$. After ϕ starts to move, one can use two extreme approximations to study the dilution of ϕ: the slow-roll approximation and the virial motion. With the slow-roll approximation, the solution (4.3) above dilutes as $\phi^4 \propto R^{-3} \propto H^2$. The opposite limit of the virial (oscillatory) motion with ϕ^6 potential gives $\dot{\phi}^2/2 = 3V$. Then the equation of motion $\ddot{\phi} + 3H\dot{\phi} + V' = 0$ can be rewritten in terms of the energy density $\rho_\phi = \dot{\phi}^2/2 + V$ as $\dot{\rho_\phi} = -3H\dot{\phi}^2 = -\frac{9}{2}H\rho_\phi$, and hence $\rho_\phi \propto R^{-9/2}$. Therefore we again obtain $\phi^4 \propto R^{-3} \propto H^2$. The true motion is somewhere between the slow-roll and virial limits but the agreement of both extremes suggests that this is the true power law of dilution. Then the amplitude ϕ always traces $\sqrt{MH}$. On the other hand, once $|\phi| \lesssim \phi_{LG}$, ϕ has a quadratic potential. In this case, the particle picture is valid for ϕ and ϕ dilutes as $|\phi|^2 \propto R^{-3}$.

For a low reheating temperature ($T_{RH} \lesssim 10^{10}$ GeV), $|\phi| \sim \phi_{LG}$ is realized in the inflaton-dominated Universe. Because ϕ tracks $\sqrt{MH}$, this occurs when $H \sim H_{LG} = m$. This corresponds to the inflaton energy density

$$\rho_\psi(t_{LG}) \sim 3m^2 M_*^2 \,, \tag{4.4}$$

where $M_* = M_{\rm Planck}/\sqrt{8\pi} = 2.4 \times 10^{18}$ GeV is the reduced Planck mass.

Lepton number density at this time is easily estimated. Define the ϕ-number density $n_\phi \equiv i(\dot{\phi}^* \phi - \phi^* \dot{\phi})$, which is related to the lepton number density as $n_L = \frac{1}{2} n_\phi$, then its equation of motion is given by

$$\dot{n}_\phi + 3H n_\phi = i\left(\frac{\partial V(\phi)}{\partial \phi} - \text{h.c.}\right) = \frac{A}{M}\,\text{Im}(\phi^4)\,. \tag{4.5}$$

As a result, we find that

$$\frac{\partial}{\partial t}(R^3 n_\phi) = R^3 \frac{A}{M}\,\text{Im}(\phi^4)\,, \tag{4.6}$$

and hence

$$(R^3 n_\phi)(t_{LG}) = \int_0^\infty dt\, R^3 \frac{A}{M}\,\text{Im}(\phi^4)\,. \tag{4.7}$$

For $t < t_{LG}$ in the integral, ϕ^6 dominates the potential and $\phi^4 \propto R^{-3}$. The integral for this time range just gives the time interval t_{LG}. For $t > t_{LG}$, ϕ^2 dominates the potential and $\phi^4 \propto R^{-6}$. The integral for this time range gives the same contribution as the previous one. Therefore the resulting lepton asymmetry is dominated by the late time contribution at $H \sim m$ and rather insensitive to the history before then. For matter-dominated Universe, $t_{LG} = \frac{2}{3} H_{LG}^{-1}$, and hence

$$n_\phi(t_{LG}) \sim \frac{4A}{3mM}\,\text{Im}(\phi^4)\,. \tag{4.8}$$

By substituting $\phi^4 \sim m^2 M^2 e^{4i\theta}$ and $A \sim m$,

$$n_\phi(t_{LG}) \sim \frac{4}{3} m^2 M \sin 4\theta\,. \tag{4.9}$$

The size of the resulting lepton number is determined by the initial phase of the field θ. This is because the potential has "valleys" along $\theta = \pi/4, 3\pi/4, 5\pi/4, 7\pi/4$ and "ridges" along $\theta = 0, \pi/2, \pi, 3\pi/2$ and the initial phase generates a rotational motion of ϕ on the complex plane when it rolls down the "slope" from the "ridge." In this sense, the lepton number is determined by a "spontaneous" CP violation because of the random phase the field acquired during the inflation.

When the inflaton decays, the Universe is reheated. With the instantaneous decay approximation, $\rho_\psi(t_{RH}) = \frac{\pi^2}{30} g_* T_{RH}^4$, and we obtain

$$\left(\frac{R_{RH}}{R_{LG}}\right)^3 = \frac{\rho_\psi(t_{LG})}{\rho_\psi(t_{RH})} \sim \frac{3m^2 M_*^2}{\frac{\pi^2}{30} g_* T_{RH}^4} . \tag{4.10}$$

Therefore, the ϕ number density at the reheating is given by

$$n_\phi(t_{RH}) = n_\phi(t_{LG}) \left(\frac{R_{LG}}{R_{RH}}\right)^3 \sim \frac{4M}{9M_*^2} \sin 4\theta \times \frac{\pi^2}{30} g_* T_{RH}^4 , \tag{4.11}$$

or using $s(t_{RH}) = \frac{2\pi^2}{45} g_* T_{RH}^3$, we obtain

$$\frac{n_\phi}{s} \sim \frac{M T_{RH}}{3M_*^2} \sin 4\theta . \tag{4.12}$$

The above estimate is valid as long as $|\phi| \sim \phi_{LG}$ is realized in the inflaton-dominated Universe. This is the case if $3m^2 M_*^2 \gtrsim \frac{\pi^2}{30} g_* T_{RH}^4$, or equivalently, $T_{RH} \lesssim 10^{10}$ GeV for $m \sim 100\,\mathrm{GeV}$. For higher reheating temperature, however, $|\phi| \sim \phi_{LG}$ happens after the reheating. At the end of the inflation, $H = H_{\rm inf}$, and the energy density of the oscillating inflaton ψ is $\rho_\psi = 3H_{\rm inf}^2 M_*^2$. Using $\rho_\psi(t_{RH}) = \frac{\pi^2}{30} g_* T_{RH}^4$, we find $(R_{\rm inf}/R_{RH})^3 \sim (\frac{\pi^2}{30} T_{RH}^4)/(3H_{\rm inf}^2 M_*^2)$. At the same time, the flat direction is diluted as

$$|\phi|^4(t_{RH}) \sim M^2 H_{\rm inf}^2 \left(\frac{R_{\rm inf}}{R_{RH}}\right)^3 \sim \frac{M^2}{3M_*^2} \times \frac{\pi^2}{30} g_* T_{RH}^4 , \tag{4.13}$$

which is still larger than ϕ_{LG}^4 for $T_{RH} \gtrsim 10^{10}$ GeV.

At the time when $|\phi|^4 \sim m^2 M^2$, the lepton number gets fixed. The scale factor at this point is determined by

$$\left(\frac{R_{RH}}{R_{LG}}\right)^3 = \left(\frac{\phi_{LG}}{\phi_{RH}}\right)^4 \sim \frac{3m^2 M_*^2}{\frac{\pi^2}{30} g_* T_{RH}^4} . \tag{4.14}$$

The radiation energy density at $t = t_{LG}$ is therefore given by

$$\rho_{rad}(t_{LG}) = \left(\frac{\pi^2}{30} g_* T_{RH}^4\right) \left(\frac{R_{RH}}{R_{LG}}\right)^4 \sim 3m^2 M_*^2 \left(\frac{3m^2 M_*^2}{\frac{\pi^2}{30} g_* T_{RH}^4}\right)^{1/3} . \tag{4.15}$$

This gives the expansion rate at this point

$$H_{LG} \sim m \left(\frac{3m^2 M_*^2}{\frac{\pi^2}{30} g_* T_{RH}^4}\right)^{1/6} . \tag{4.16}$$

The evolution of n_ϕ is solved as in the previous case. For radiation-dominated Universe, $t_{LG} = \frac{1}{2} H_{LG}$, and hence eq. (4.7) gives

$$n_\phi(t_{LG}) \sim \frac{1}{2} m^2 M \left(\frac{3m^2 M_*^2}{\frac{\pi^2}{30} g_* T_{RH}^4}\right)^{-1/6} \sin 4\theta , \tag{4.17}$$

where we used $\phi^4 \sim m^2 M^2 e^{4i\theta}$ and $A \sim m$. If we naively adopt the above relation, n_ϕ is bigger than $m^2 M \sim m|\phi_{LG}|^2$ for $\theta \sim O(1)$. However, at this point, the ϕ motion allows a particle interpretation because the potential is dominated by the quadratic piece, and the maximum possible number of particles is given by $\sim m|\phi|^2$. If we take the above estimate literally, n_ϕ may be bigger than $m|\phi|^2$: a contradiction. Then the conclusion is that n_ϕ has already saturated before this stage. Therefore, some of the previous discussions for $T_{RH} \gtrsim (mM_*)^{1/2}$ are somewhat irrelevant, and we can simply estimate n_ϕ at $\phi \sim \phi_{LG}$ to be $\sim m^2 M$. Then, using the entropy at t_{LG}:

$$s(t_{LG}) = s(t_{RH}) \left(\frac{R_{RH}}{R_{LG}} \right)^3 \sim \frac{4m^2 M_*^2}{T_{RH}}, \tag{4.18}$$

we obtain a similar ϕ-to-entropy ratio as given in eq. (4.12) up to an $O(1)$ coefficient. We use the result given in eq. (4.12) for the numerical estimation below.

With the estimated n_ϕ/s (4.12), the lepton-to-entropy ratio is given by

$$Y_L \equiv \frac{n_L}{s} = \frac{n_\phi}{2s} \sim \frac{1}{6} \frac{M T_{RH}}{M_*^2} \sin 4\theta = \frac{1}{6} \frac{v_u^2 T_{RH}}{m_\nu M_*^2} \sin 4\theta \,. \tag{4.19}$$

Since the lepton number is inversely proportional to the mass of the neutrino, it is dominated by the lightest mass eigenvalue m_{ν_1}, possibly mostly the electron neutrino state.

So far the discussion used only the flat-space potential (case (1)). For case (2), the discussion is modified only slightly. The field amplitude dilutes as $\phi \sim \sqrt{MH}$ because the negative mass squared $-H^2/16\pi^2$ is smaller than the expansion rate and hence does not modify the discussion with flat-space potential. The main issue here is if the phase direction of the field gets settled to the "valley." Because of loop-suppressed effect of the expansion in the potential, the potential along the phase direction is given by $((A + cH/16\pi^2)\phi^4 + \text{c.c.})/8M$ with $c = O(1)$, and the curvature along the phase direction is of the order of $cH\phi^2/16\pi^2 M \sim cH^2/16\pi^2$. This is always smaller than the expansion rate and hence the motion along the phase direction is critically damped. Therefore the initial phase of the potential is kept intact until the expansion slows down to $H \sim m$ when the flat-space potential is recovered and the field starts rolling down the slope from the "ridge". The estimate of the resulting lepton asymmetry is hence the same as in the case (1) with the flat-space potential.

The discussion is different for case (3) where the effect of the expansion introduces a large supersymmetry breaking effect of $O(H)$ in the potential. This case was discussed in detail in [21], and we summarize the situation briefly here. The field amplitude always tracks the minimum of the potential $\phi \sim \sqrt{MH}$ because the curvature $-H^2|\phi|^2$ is of the same order of magnitude as the expansion rate itself and field always rolls down to the minimum at any given time. Therefore the field amplitude is the same as other two cases even though the reason for it is quite different. Similarly, the potential along the phase direction $((A + cH)\phi^4 + \text{c.c.})/8M$

with $c = O(1)$ produces the curvature along the phase direction is also of the order of $cH\phi^2/M \sim H^2$ and hence the field settles to the "valley" and tracks it. However, the potential can generate a rotational motion when H becomes smaller than A because of a possible relative phase between two terms A and cH. The phase for the "valleys" quickly shifts to that of the flat-space potential at this moment and the field which tracked the valley with the H effect starts moving to the new valley with the A term only. The estimate of the rotational motion is basically the same as the earlier two cases and hence the resulting lepton asymmetry as well. What is different from the previous cases is that the CP violation is in the relative phase between A and cH and originates from the microscopic physics rather than a spontaneous random phase.

To summarize, in all cases (1), (2), (3) considered, the final lepton asymmetry from the motion of the D-flat direction $\tilde{L} = H_u$ is given by eq. (4.19) which confirms the original estimate in [15]. Therefore the result is quite robust and model independent.

5. Estimate of baryon asymmetry

In order to proceed to a numerical estimate of the generated lepton asymmetry, we need to have some idea on the lightest neutrino mass from the currently available data. Here we have to rely on certain assumptions.

We assume hierarchical neutrino masses similar to quark and charged lepton masses, and hence that Δm^2 from neutrino oscillation give larger of the mass eigenvalues between two neutrino states relevant to the particular oscillation mode. The atmospheric neutrino data require that the largest two eigenvalues have $\Delta m^2 = 10^{-3}$–$10^{-2}\,\mathrm{eV}^2$, and hence according to our assumption, $m_{\nu_3} = 0.03$–$0.1\,\mathrm{eV}$. The small angle MSW solution suggests $\Delta m^2 \simeq 4$–$10\times 10^{-6}\,\mathrm{eV}^2$, which should give the second lightest mass eigenvalue $m_{\nu_2} \simeq 0.0020$–$0.0032\,\mathrm{eV}$ and the mixing angle $\sin^2 2\theta_{MNS} \simeq 1\times 10^{-3}$–$1 \times 10^{-2}$ [22]. The estimate of the lightest mass eigenvalue depends on the assumptions on the mass matrix. If the MNS (Maki-Nakagawa-Sakata) mixing angle [23] is solely due to the neutrino Majorana matrix, the lightest eigenvalue is likely to be around $m_{\nu_1} \sim m_{\nu_2}\theta^2_{MNS} \sim (0.5$–$8) \times 10^{-6}\,\mathrm{eV}$. This approximate relation between the mass eigenvalue and the angle is quite generic unless most of the mixing is attributed to the charged lepton sector. For instance the models based on coset-space family unification on $E_7/\,\mathrm{SU}(5)\times \mathrm{U}(1)^3$ [24], string-inspired anomalous U(1) [25], and SU(5) grand unified model with abelian flavor symmetry [26] all give this approximate relation. However if the mixing angle is coming from the charged lepton sector, the mass eigenvalue can be smaller.

For the large angle MSW solution, $m_{\nu_2} \simeq 0.003$–$0.01\,\mathrm{eV}$, and the lightest eigenvalue is likely to be not much smaller than m_{ν_2} if the large MNS angle comes from the neutrino mass matrix, but can be much smaller if the MNS angle originates in the charged lepton mass matrix. This might be considered a fine-tuning however because the determinant of the mass matrix (either charged lepton or neutrino) should be much smaller than the typical size of the elements. The vacuum solution most likely gives $m_{\nu_1} \simeq m_{\nu_2} \simeq 10^{-5}$ eV, which may also need a fine-tuning in the mass matrices. For the purpose of the estimation below, we take $m_{\nu_1} \sim 2.5 \times 10^{-6}$ eV, preferred by the small angle MSW solution, but the result can be easily scaled according to the value of m_{ν_1}.

The reheating temperature cannot be arbitrarily high because of the cosmological constraints on the gravitino production. For the range of gravitino mass usually considered, $m_{3/2} \simeq 500$ GeV or so, the primordial gravitinos will decay with the lifetime of $\tau_{3/2} \simeq 2 \times 10^5$ sec and therefore can potentially destroy the light elements synthesized by Big-Bang Nucleosynthesis. The most robust constraint is from the thermally produced gravitino, which requires the reheating temperature to be less than $T_{RH} \lesssim 10^9$ GeV [11]. Another constraint is obtained from the non-thermal production after the inflationary epoch, which may give a more stringent constraint on the reheating temperature [27]. However, the number density of the non-thermally produced gravitino is model-dependent. For example, for the chaotic inflation model with the inflaton superpotential of $W_{\text{inflaton}} \sim \frac{1}{2} M_{\text{inflaton}} \phi^2_{\text{inflaton}}$, the gravitino to entropy ratio is of order $M_{\text{inflaton}} T_{RH}/M_*^2$, which is smaller than that of the thermally produced gravitino. Because of the model-dependence, we do not include the constraint from the non-thermally produced gravitino into our discussion, and we take the canonical value of $T_{RH} = 3 \times 10^8$ GeV for the estimation below.[6]

Now come the numerical estimates. We find the lepton asymmetry (4.19) to be

$$Y_L \sim 1.1 \times 10^{-10} \left(\frac{2.5 \times 10^{-6} \text{ eV}}{m_{\nu_1}} \right) \left(\frac{T_{RH}}{3 \times 10^8 \text{ GeV}} \right) \sin 4\theta \,. \tag{5.1}$$

From the chemical equilibrium between lepton and baryon numbers, we obtain $Y_B = 0.35 Y_L$ [28], while the number density of photons now is related to the entropy density by $n_\gamma / s = (410.89 \text{ cm}^{-3})/(2892.4 \text{ cm}^{-3})$, and in the end we find

$$\eta = \frac{n_B}{n_\gamma} \simeq 2.46 Y_L \sim 2.6 \times 10^{-10} \left(\frac{2.5 \times 10^{-6} \text{ eV}}{m_{\nu_1}} \right) \left(\frac{T_{RH}}{3 \times 10^8 \text{ GeV}} \right) \sin 4\theta \,. \tag{5.2}$$

Recall that we used the preferred range by the small angle MSW solution $m_{\nu_1} \sim (0.5$–$8) \times 10^{-6}$ eV. This estimate should be compared to the value determined from the Big-Bang Nucleosynthesis $\eta \simeq (4-6) \times 10^{-10}$. Therefore this mechanism can generate

[6]Moduli fields may also be dangerous. In particular, if they have large initial amplitudes of $O(M_*)$, their primordial number density becomes so large that the Big-Bang Nucleosynthesis is seriously damaged. The initial amplitude may be suppressed, for example, if the minimum of the moduli potential is given by a symmetry enhanced point [18]. Non-thermal production of the moduli may be also important, but is also model-dependent [27]. In particular, if the modulus have a large effective mass of $O(H_{\text{inf}})$ during the inflation, or if the modulus is conformally coupled, then its non-thermal production after the inflation is not effective.

the cosmic baryon asymmetry of the correct order of magnitude without conflicting the constraints on the cosmological gravitino production within the range of neutrino mass preferred by the atmospheric neutrino data and the small angle MSW solution to the solar neutrino problem. However, one should note that such a small neutrino mass can be accommodated even with other solutions to the solar neutrino problem (i.e. large angle MSW and vacuum oscillation solutions) by attributing the MNS angle to the charged lepton sector or by a fine tuning in the neutrino mass matrix.

6. Conclusion

To conclude, we have shown that the minimal supersymmetric standard model together with Majorana mass of neutrinos can generate the cosmic baryon asymmetry without any additional new physics, hence the minimal supersymmetric leptogenesis. The required neutrino mass is consistent with the mass matrices proposed to explain the atmospheric as well as solar neutrino data with the small angle MSW solution. Other solutions to the solar neutrino problem can also be accommodated if the lightest neutrino mass eigenvalue is of the order of 10^{-6} eV. The reheating temperature can be low enough to avoid cosmological gravitino problem.

Acknowledgments

We thank Steve Martin and Pierre Ramond for helpful conversations, and Gian Giudice and Antonio Riotto for useful comments. H.M. also thanks Masahiro Kawasaki and Tsutomu Yanagida for encouraging him to write up this work. We would like to thank Aspen Center for Physics where this work was completed. This work was supported in part by the U.S. Department of Energy under Contracts DE-AC03-76SF00098, and in part by the National Science Foundation under grants PHY-95-13835 and PHY-95-14797. T.M. was also supported by the Marvin L. Goldberger Membership. H.M. was also supported by Alfred P. Sloan Foundation.

References

[1] Super-Kamiokande collaboration, *Current status of nucleon decay searches with* Super-Kamiokande, hep-ex/9903029.

[2] Super-Kamiokande collaboration, *Evidence for oscillation of atmospheric neutrinos*, *Phys. Rev. Lett.* **81** (1998) 1562 [hep-ex/9807003].

[3] Soudan-2 collaboration, *The atmospheric neutrino flavor ratio from a 3.9 fiducial kiloton-year exposure of* Soudan 2, *Phys. Lett.* **B 449** (1999) 137 [hep-ex/9901024].

[4] MACRO collaboration, *Measurement of the atmospheric neutrino-induced upgoing muon flux using* Macro, *Phys. Lett.* **B 434** (1998) 451 [hep-ex/9807005].

[5] T. Yanagida, in *Proceedings of the workshop on unified theory and baryon number of the universe*, O. Sawada and A. Sugamoto eds., KEK, 1979, p. 95;
M. Gell-Mann, P. Ramond and R. Slansky, in *Supergravity*, P. van Niewwenhuizen and D. Freedman eds., North Holland, Amsterdam, 1979.

[6] N. Arkani-Hamed, S. Dimopoulos, G. Dvali and J. March-Russell, *Neutrino masses from large extra dimensions*, `hep-ph/9811448`;
K.R. Dienes, E. Dudas and T. Gherghetta, *Light neutrinos without heavy mass scales: a higher-dimensional seesaw mechanism*, *Nucl. Phys.* **B 557** (1999) 25 [`hep-ph/9811428`].

[7] M. Fukugita and T. Yanagida, *Baryogenesis without grand unification*, *Phys. Lett.* **B 174** (1986) 45.

[8] V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, *On the anomalous electroweak baryon number nonconservation in the early universe*, *Phys. Lett.* **B 155** (1985) 36.

[9] W. Buchmüller and T. Yanagida, *Quark lepton mass hierarchies and the baryon asymmetry*, *Phys. Lett.* **B 445** (1999) 399 [`hep-ph/9810308`].

[10] L. Randall and R. Sundrum, *Out of this world supersymmetry breaking*, *Nucl. Phys.* **B 557** (1999) 79 [`hep-th/9810155`];
G.F. Giudice, M.A. Luty, H. Murayama and R. Rattazzi, *Gaugino mass without singlets*, *J. High Energy Phys.* **12** (1998) 027 [`hep-ph/9810442`].

[11] E. Holtmann, M. Kawasaki, K. Kohri and T. Moroi, *Radiative decay of a long-lived particle and big-bang nucleosynthesis*, *Phys. Rev.* **D 60** (1999) 023506 [`hep-ph/9805405`].

[12] T. Asaka, K. Hamaguchi, M. Kawasaki and T. Yanagida, *Leptogenesis in inflaton decay*, *Phys. Lett.* **B 464** (1999) 12 [`hep-ph/9906366`].

[13] G.F. Giudice, M. Peloso, A. Riotto and I. Tkachev, *Production of massive fermions at preheating and leptogenesis*, *J. High Energy Phys.* **08** (1999) 014 [`hep-ph/9905242`].

[14] I. Affleck and M. Dine, *A new mechanism for baryogenesis*, *Nucl. Phys.* **B 249** (1985) 361.

[15] H. Murayama and T. Yanagida, *Leptogenesis in supersymmetric standard model with righthanded neutrino*, *Phys. Lett.* **B 322** (1994) 349 [`hep-ph/9310297`].

[16] D.S. Salopek, *Consequences of the cobe satellite for the inflationary scenario*, *Phys. Rev. Lett.* **69** (1992) 3602.

[17] A.D. Linde, *Chaotic inflation*, *Phys. Lett.* **B 129** (1983) 177; *The new mechanism of baryogenesis and the inflationary universe*, *Phys. Lett.* **B 160** (1985) 243.

[18] M. Dine, L. Randall and S. Thomas, *Supersymmetry breaking in the early universe*, *Phys. Rev. Lett.* **75** (1995) 398 [`hep-ph/9503303`].

[19] M.K. Gaillard, H. Murayama and K.A. Olive, *Preserving flat directions during inflation*, *Phys. Lett.* **B 355** (1995) 71 [hep-ph/9504307].

[20] See, for a review, A. Linde, *Particle physics and inflationary cosmology*, Harwood, Chur, Switzerland, 1990, Contemporary concepts in physics, 5.

[21] M. Dine, L. Randall and S. Thomas, *Baryogenesis from flat directions of the supersymmetric standard model*, *Nucl. Phys.* **B 458** (1996) 291 [hep-ph/9507453].

[22] For a recent analysis, see M.C. Gonzalez-Garcia, P.C. de Holanda, C. Pena-Garay and J.W.F. Valle, *Status of the msw solutions of the solar neutrino problem*, *Nucl. Phys.* **B 573** (2000) 3 [hep-ph/9906469].

[23] Z. Maki, M. Nakagawa, and S. Sakata, *Remarks on the unified model of elementary particles*, *Prog. Theor. Phys.* **28** (1962) 870.

[24] J. Sato and T. Yanagida, *Large lepton mixing in a coset-space family unification on* $E_7/\mathrm{SU}(5)\times\mathrm{U}(1)^3$, *Phys. Lett.* **B 430** (1998) 127 [hep-ph/9710516].

[25] N. Irges, S. Lavignac and P. Ramond, *Predictions from an anomalous* U(1) *model of Yukawa hierarchies*, *Phys. Rev.* **D 58** (1998) 035003 [hep-ph/9802334];
J.K. Elwood, N. Irges and P. Ramond, *Family symmetry and neutrino mixing*, *Phys. Rev. Lett.* **81** (1998) 5064 [hep-ph/9807228].

[26] G. Altarelli and F. Feruglio, *A simple grand unification view of neutrino mixing and fermion mass matrices*, *Phys. Lett.* **B 451** (1999) 388 [hep-ph/9812475].

[27] G.F. Giudice, I. Tkachev and A. Riotto, *Non-thermal production of dangerous relics in the early universe*, *J. High Energy Phys.* **08** (1999) 009 [hep-ph/9907510].

[28] S.Y. Khlebnikov and M.E. Shaposhnikov, *The statistical theory of anomalous fermion number nonconservation*, *Nucl. Phys.* **B 308** (1988) 885.