We investigate thermal leptogenesis in a supersymmetric neutrinophilic Higgs model by taking phenomenological constraints into account, where, in addition to the minimal supersymmetric standard model, we introduce an extra Higgs field with a tiny vacuum expectation value (VEV) which generates neutrino masses. Thanks to this tiny VEV of the neutrinophilic Higgs, our model allows to reduce the mass of the lightest right-handed (s)neutrino to be $O(10^5)$ GeV as keeping sufficiently large CP asymmetry in its decay. Therefore, the reheating temperature after inflation is not necessarily high, hence this scenario is free from gravitino problem.

The origin of cosmological baryon asymmetry is one of the most important questions in both particle physics and cosmology. Among various mechanisms of generating the suitable baryon asymmetry, leptogenesis [1] is one of the most attractive scenarios. Particularly, thermal leptogenesis requires only thermal excitation of right-handed Majorana neutrinos which generate tiny neutrino masses via a seesaw mechanism [2], and provides several implications for the spectrum [3] of light neutrino masses confirmed by neutrino oscillation experiments [4, 5]. However, a realization of thermal leptogenesis has a difficulty of “gravitino problem” [6] in supersymmetric models with R-parity. In order to avoid the overproduction of gravitinos, the reheating temperature after inflation $T_R$ must not be so high to thermalize right-handed (s)neutrinos [7]. Therefore, gravitino problem is a serious obstacle in a usual Type-I seesaw [2], where tiny neutrino masses of order $0.1$ eV is obtained through superheavy right-handed neutrinos.

How about an alternative idea, a neutrinophilic Higgs doublet model [8–13]? Here the smallness of neutrino masses originates from a tiny VEV of neutrinophilic Higgs doublet, and neutrino Yukawa couplings are not tiny anymore. Recently, we have shown that thermal leptogenesis could work at a low energy scale in a neutrinophilic Higgs doublet model without gravitino problem [14]. However, it is also worried that enlarge neutrino Yukawa couplings might give rise to sizable processes of lepton flavor violations (LFVs). Thus, in this paper, we will show that thermal leptogenesis surely works without gravitino problem in a supersymmetric neutrinophilic Higgs doublet model after carefully taking other phenomenological constraints into account.

The supersymmetric neutrinophilic Higgs model has a pair of neutrinophilic Higgs doublets $H_\nu$ and $H_{\nu'}$ in addition to up- and down-type two Higgs doublets $H_u$ and $H_d$ in the minimal supersymmetric standard model (MSSM). A discrete $Z_2$-parity to discriminate $H_u(H_d)$ from $H_\nu(H_{\nu'})$ is also introduced, and its charges (and also lepton number) are assigned as the following table. Under the discrete symmetry, the superpotential is given by

| fields                        | $Z_2$-parity | lepton number |
|-------------------------------|--------------|---------------|
| MSSM Higgs doublets, $H_u, H_d$ | +            | 0             |
| new Higgs doublets, $H_{\nu}, H_{\nu'}$ | −            | 0             |
| right-handed neutrinos, $N$    | −            | 1             |
| others                        | +±1: leptons, 0: quarks |               |

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\[ W = y^u \bar{Q} H_u U_R + y^d \bar{Q} H_d D_R + y^l \bar{L} H_d E_R \\
+ y^\nu \bar{L} H_\nu N + \frac{1}{2} MN^2 \\
+ \mu H_u H_d + \mu' H_\nu H_\nu + \rho H_u H_\nu + \rho' H_\nu H_d, \quad (1) \]

where we omit generation indexes. The Z_2-parity plays a crucial role of suppressing tree-level flavor changing neutral currents (FCNCs), and is assumed to be softly broken by tiny parameters of \( \rho \) and \( \rho' (\ll \mu, \mu') \). We expect that supersymmetry breaking soft squared masses can trigger suitable electro-weak symmetry breaking. The Higgs potential is given by

\[ V = |\mu|^2 (H_u^\dagger H_u + H_d^\dagger H_d) + |\mu'|^2 (H_\nu^\dagger H_\nu + H_\nu'^\dagger H_\nu') \\
+ \frac{\tan^2 \beta - 1}{2} \sum_a \frac{g_a^2}{2} \left( H_u^\dagger \tau^a H_u + H_d^\dagger \tau^a H_d + H_\nu^\dagger \tau^a H_\nu + H_\nu'^\dagger \tau^a H_\nu' \right)^2 \\
+ m^2_{H_u} H_u^\dagger H_u + m^2_{H_d} H_d^\dagger H_d + m^2_{H_\nu} H_\nu^\dagger H_\nu + m^2_{H_\nu'} H_\nu'^\dagger H_\nu' \\
+ B \mu H_u \cdot H_d + B' \mu' H_\nu \cdot H_\nu' + B \rho H_u \cdot H_\nu + B' \rho' H_\nu' \cdot H_d + h.c., \]

where we have omitted tiny \( \rho^2 \) and \( \rho'^2 \) mass terms. \( \tau^a \) and dot represent a generator of SU(2) and its anti-symmetric product, respectively, and \( g_1 (g_2) \) is a gauge coupling constant of U(1)_Y (SU(2)_L). \( m^2_{H_u}, m^2_{H_d}, m^2_{H_\nu}, m^2_{H_\nu'} \) and \( B, B', B'' \) are soft SUSY breaking parameters. The tiny soft Z_2-breaking parameters, \( \rho, \rho' \), generate a large hierarchy of \( v_{u,d} (\equiv (H_{u,d}) \gg v_{\nu,\nu'} (\equiv (H_{\nu,\nu'})) \) through stationary conditions,

\[ \left( \begin{array}{cc} m^2_{H_u} + \mu^2 + m^2_{H_d} \tan^2 \beta - 1 & -B \mu' \mu' \\
-B \rho' \mu' & m^2_{H_\nu} + \mu'^2 - m^2_{H_\nu'} \tan^2 \beta - 1 \end{array} \right) \left( \begin{array}{c} v_\nu \\
v_{\nu'} \end{array} \right) \approx \left( \begin{array}{cc} -\rho' \rho' & \rho' \rho' \\
-\rho \rho & -\rho \rho \end{array} \right) \left( \begin{array}{c} v_u \\
v_d \end{array} \right), \quad (3) \]

For example, \( v_\nu \sim 1 \text{ GeV} \) is obtained from \( \rho, \rho' \sim 1 \text{ GeV} \) with Higgs mass parameters of \( O(10^2) \) GeV. At the vacuum of \( v_{\nu,\nu'} \ll v_{u,d} \) that we are interested in, physical Higgs bosons originated from \( H_{u,d} \) are almost decoupled from those from \( H_{\nu,\nu'} \). The former, \( H_{u,d} \), almost constitute Higgs bosons in the MSSM; two CP-even Higgs boson \( h \) and \( H \), one CP-odd Higgs boson \( A \), and charged Higgs boson \( H^\pm \), while the latter, \( H_{\nu,\nu'} \), constitute two CP-odd Higgs bosons \( H_{2,3} \), two CP-odd bosons \( A_{2,3} \), and two charged Higgs bosons \( H^+_{2,3} \). The two physical charged Higgs bosons are given by

\[ \left( \begin{array}{c} H^+_{2,3} \\
H^0_{2,3} \end{array} \right) = \left( \begin{array}{cc} \cos \alpha_c & -\sin \alpha_c \\
\sin \alpha_c & \cos \alpha_c \end{array} \right) \left( \begin{array}{c} H^+_3 \\
H^0_3 \end{array} \right), \quad (4) \]

where \( \tan 2 \alpha_c = 2 B \mu' / (m^2_{H_u} - m^2_{H_\nu} + (m^2_\nu - 2m^2_{\nu'} \tan^2 \beta - 1) \) and \( \tan \beta = v_u / v_d \).

Through the seesaw mechanism, masses of light neutrinos are given by

\[ m_{ij} = \sum_k \frac{y^\nu_k v_\nu y^{\nu T}_k y^\nu_j}{M_k}. \quad (5) \]

For fixed right-handed neutrino masses, a tiny VEV of \( v_\nu \) requires larger neutrino Yukawa couplings \( y^\nu \) than conventional seesaw scenarios. The neutrino masses may be also received radiative corrections as \( m^\text{loop}_{ij} \sim -\lambda v_\nu^2 m^2_{\tilde{H}^\pm_{2,3}} /(8\pi^2 M) \), where we assume \( M \gg m^2_{\tilde{H}^\pm_{2,3}} \) and \( \lambda \) is a coupling of one-loop induced scalar interaction, \( \lambda (H_\nu \cdot H_\nu) \). Notice that the radiative induced mass is smaller than the tree-level mass of Eq. (5) as long as \( \lambda < 16 \pi^2 v_\nu^2 / v^2 \). Thus, we can neglect radiative corrections of neutrino masses, since our model induces \( \lambda \sim g^2 \rho^2 / (32 \pi^2 m^2_{\tilde{H}^\pm_{2,3}}) \sim 10^{-10} \), where \( m_{\tilde{H}^\pm} \) is a chargino mass. For this estimation, we have used \( v_{\nu}/v_d \sim 10^{-2} \) which will be a suitable parameter region in the following discussions. Actually, there are two 1-loop diagrams which contribute \( \lambda \), but the chargino 1-loop diagram dominates a Higgs 1-loop diagram, and we neglect the latter. Notice that (s)top 1-loop diagram is negligible due to the Z_2-parity. Anyhow, the value of \( \lambda \) is tiny, since it is not induced until 1-loop diagrams including both Z_2- and SUSY-breaking effects. The neutrino mass matrix in Eq. (5) can reproduce neutrino oscillation experiments, and we will use a concrete value of \( \Delta m^2_{\text{atm}} \) in the following analyses.
Now, let us discuss thermal leptogenesis in this model. A resultant baryon asymmetry generated via thermal leptogenesis is generally given by

\[ \frac{n_b}{s} \simeq C \nu \frac{\varepsilon}{g_*}, \]  

(6)

where \( g_*|_{T=M_1} = \mathcal{O}(100) \) is the effective degrees of freedom of relativistic particles in thermal bath, and \( \varepsilon \) is the total CP asymmetry of right-handed (s)neutrino decay. Dilution (or efficiency) factor \( \kappa \leq \mathcal{O}(0.1) \) denotes the dilution by washout processes, and the coefficient \( C \) is a factor of the conversion from lepton to baryon asymmetry by the sphaleron [22]. The decay rate of the lightest right-handed neutrino \( N_1 \) into left-handed (s)leptons and \( H_\nu(\tilde{H}_\nu) \)-like Higgs bosons (higgsinos) \( \Gamma_{N_1} \) and that of the lightest right-handed sneutrino \( \tilde{N}_1 \) into left-handed sleptons (leptons) and \( H_\nu(\tilde{H}_\nu) \)-like Higgs bosons (higgsinos) \( \Gamma_{\tilde{N}_1} \) are given by

\[ \Gamma_{N_1} = \Gamma_{\tilde{N}_1} = \sum_j \frac{y_{\nu j}^2 y_{i j}^2}{4\pi} |M_1| = \frac{(y_{\nu}^\dagger y_{\nu})_{11}^2}{4\pi} |M_1|, \]  

(7)

We would note here that physical mass eigenstate of \( H_\nu(\tilde{H}_\nu) \)-like Higgs boson (higgsino) has tiny component of \( H_\nu(\tilde{H}_\nu) \) through the tiny \( Z_2 \)-breaking parameter \( \eta \). The condition for out of equilibrium in decay of right-handed (s)neutrino \( \Gamma_{N_1(\tilde{N}_1)} < \mathcal{O}(10^3) \) requires that the lightest left-handed neutrino is almost massless \( m_1 \simeq 0 \) and \( y_{\nu j}^2 \) are very small, where \( H \) is the Hubble parameter and \( T \) is the temperature of radiation. For the neutrino Yukawa couplings of \( y_{\nu 1}^2 \ll y_{\nu 2}^2, y_{\nu 3}^2 \) and hierarchical right-handed neutrino mass spectrum [18], the total CP asymmetry of right-handed (s)neutrino decay is given by

\[ \varepsilon \equiv \varepsilon(N \to lH) + \varepsilon(N \to \bar{L}H) + \varepsilon(\tilde{N} \to lH) + \varepsilon(\tilde{N} \to \bar{L}H) \]

\[ \simeq -\frac{3}{16\pi} 10^{-5} \left( \frac{0.1\text{GeV}}{v_{\nu}} \right)^2 \left( \frac{M_1}{10^3\text{GeV}} \right) \left( \frac{m_\nu}{0.05\text{eV}} \right) \sin\theta. \]  

(8)

Here \( \theta \) is an effective CP violating phase, which is significantly enhanced due to the tiny \( v_\nu \). In order to obtain the observed baryon asymmetry in our Universe \( n_b/s \simeq 10^{-10} \) [19], \( \varepsilon \gtrsim 10^{-7} \) is required. For the conventional Type-I seesaw in the MSSM with superheavy right-handed neutrinos (where the neutrino Dirac mass term is generated through \( v_u \)), \( \varepsilon \gtrsim 10^{-7} \) means \( M_1 \gtrsim 10^9 \text{ GeV} \), which is so-called Davidson-Ibarra bound for models with hierarchical right-handed neutrino mass spectrum [20, 21]. In contrast, \( v_\nu \) is replaced by \( v_{\nu\pm}(\ll v_u) \) in our model, and as the result, an enough large \( \varepsilon \) can be obtained even for a smaller \( M_1 \) than that derived by Davidson-Ibarra bound.

Lepton number violating scatterings act as washout processes of the generated lepton number asymmetry. Those scatterings are classified into two classes; the lepton number is violated by one \( \Delta \nu = 1 \) scattering rates are proportional to \( \Gamma_{N_1} \) and hence can be minimized by its appropriate choice. We should notice that \( \Delta L = 1 \) processes such as \( LN \to H \to QA \) are negligible, since the mixing between \( H_\nu(\tilde{H}_\nu) \) and \( H_\nu(\tilde{H}_\nu) \), \( H_u(\tilde{H}_u) \) is negligible due to tiny \( Z_2 \) breaking. On the other hand, the \( \Delta L = 2 \) scatterings are potentially dangerous, and relevant scatterings for the MSSM have been studied in Ref. [22]. The decoupling condition for \( \Delta L = 2 \) lepton number violating scatterings \( \gamma_{\Delta \nu} \) in Ref. [22] is applicable to our model, which is roughly estimated as

\[ \sum_i \left( \sum_j \frac{y_{\nu j}^2 y_{i j}^2 v_{\nu j}^2}{M_j} \right)^2 < \pi^3 \zeta(3) \sqrt{\frac{\pi^2 g_*}{90}} \sqrt{\frac{\varepsilon_{\nu}}{T M_P}}, \]  

(9)

for \( T < M_1 \). Other scattering processes also give similar conditions. For a lower \( v_\nu \), washout processes are more significant. Inequality (9) gives the lower bound on \( v_\nu \) to avoid too strong washout.

As we have shown above, a sufficient CP violation \( \varepsilon = \mathcal{O}(10^{-6}) \) can be realized for \( v_\nu = \mathcal{O}(1) \) GeV in the hierarchical right-handed neutrino with \( M_1 \) of \( \mathcal{O}(10^8 - 10^9) \) GeV. This implies that the reheating temperature after inflation \( T_R \) of \( \mathcal{O}(10^9) \) GeV is high enough to produce right-handed neutrinos by thermal scatterings. Thus, this class of model with \( v_\nu = \mathcal{O}(1) \) GeV is a solution to compatible with thermal leptogenesis in gravity mediated supersymmetry breaking with unstable gravitino.

Next, let us investigate phenomenological constraints in our model. The most severe constraint comes from LFV decay processes, particularly, \( \mu \to e\gamma \). There are LFV processes from 1-loop processes triggered by \( y_{\nu}^\dagger \) loops of \( N \cdot H^\pm \)
and $\tilde{N} - \tilde{\chi}^\pm$) in addition to the MSSM processes. A branching ratio of the LFV is given by

$$B(l_\alpha \to l_\beta \gamma) = \frac{30_{\text{exp}}}{64\pi G_F} \left| \sum_i y_{i\alpha \beta} \right|^2 \left\{ \cos^2 \alpha_i \frac{M^2}{M_{H^\pm_i}^2} F \left( \frac{M^2}{M_{H^\pm_i}^2} \right) + \sin^2 \alpha_i \frac{M^2}{M_{H^\mp_i}^2} F \left( \frac{M^2}{M_{H^\mp_i}^2} \right) - \frac{M_{\tilde{\chi}^\pm}}{2m_\mu M_{\tilde{N}_i}} G \left( \frac{M_{\tilde{\chi}^\pm}^2}{M_{\tilde{N}_i}^2} \right) \right\} + \text{MSSM processes}$$

where

$$F(x) = \frac{1}{6(1-x)^4} (1 - 6x + 3x^2 + 2x^3 - 6x^2 \ln x),$$

$$G(x) = \frac{1}{(1-x)^3} (-3 + 4x - x^2 - 2 \ln x).$$

Here $G_F$ is the Fermi coupling constant. $M_{\tilde{N}_i}$ is the mass of $N_i$, and $M_{\tilde{\chi}^\pm}$ is the mass of $\tilde{H}_\nu(\tilde{H}_\nu')$-like chargino. In all parameter region except for $M_{\tilde{\chi}}^2 \approx H_{3,3}^\pm$, the chargino-loop contribution involving $G(M_{\tilde{\chi}^\pm}^2/M_{\tilde{N}_i}^2)$ is dominant and the charged Higgs boson loop contribution depending upon $F(M_{\tilde{\chi}^\pm}^2/M_{H_{3,3}^\pm}^2)$ is negligible. Thus, almost independent from the charged Higgs boson masses $H_{3,3}^\pm$, and its mixing angle $\alpha_i$, the additional contribution to the LFV decay is estimated as $B(l_\alpha \to l_\beta \gamma) \lesssim 10^{-13}$, which is much smaller than the current experimental bounds, $B(\mu \to e\gamma) < 1.2 \times 10^{-11}$, $B(\tau \to e\gamma) < 3.3 \times 10^{-8}$, and $B(\tau \to \mu\gamma) < 4.4 \times 10^{-8}$ [22]. We here take a parameter region where leptogenesis effectively works without washout effects. The similar diagram (initial and final states are both muon) induces a deviation of muon anomalous magnetic moment $a_\mu \equiv (g_\mu - 2)/2$. Similarly, additional contributions to $a_\mu$ from above loop processes is turned out to be $\Delta a_\mu = O(10^{-15})$, which is sufficiently tiny.

We here summarize all conditions for successful thermal leptogenesis, and the result is presented in the Figure 1. The horizontal axis is the VEV of neutrino Higgs $v_\nu$ and the vertical axis is the mass of the lightest right-handed neutrino $M_1$ in hierarchical right-handed neutrino mass spectrum. In the brown region, the lightest right-handed neutrino decay into $H_\nu$-like Higgs boson and lepton is kinematically not allowed. In turquoise region corresponds to inequality of Eq.(9), where $\Delta L = 2$ washout effect is too strong. The red and green line are the contours of the CP asymmetry of $\varepsilon = 10^{-6}$ and $10^{-7}$, respectively. Thus, in the parameter region near above the line of $\varepsilon = 10^{-7}$, thermal leptogenesis easily works even with hierarchical masses of right-handed neutrinos.

We have investigated thermal leptogenesis in a supersymmetric neutrinoophilic Higgs doublet model with taking account of phenomenological constraints and gravitino problem. One of the attraction of neutrinoophilic Higgs models is that the neutrino Yukawa couplings are not necessarily tiny anymore. They can enhance the CP asymmetry of right-handed (s)neutrino decay, however might also enhance the washout rate of generated lepton asymmetry and magnitudes of LFV processes, simultaneously. We have found that the suitable baryon asymmetry is reproduced with
the suitable neutrino masses of $O(10^{-1})$ eV, in which expected LFVs are consistent with current experiments and the strong $\Delta L = 2$ washout can be avoided. To generate and thermalize relatively light right-handed neutrino with mass of $O(10^5)$ GeV, the reheating temperature is low enough to avoid gravitino problem.

At the end, we comment on gauge coupling unification (GCU). It can be achieved by introducing extra vector-like $SU(3)_c$-triplet particles, $d, \bar{d}$. We also introduce an additional $Z_2$-parity, and make only $d, \bar{d}$ have odd-charge of it. Therefore, $d, \bar{d}$ have no Yukawa interactions with ordinal quarks and leptons as possessing their heavy masses of $W \sim \mu d \bar{d}$. This field content is similar to so-called Nelson-Barr model [24] and its supersymmetric version proposed in Ref. [25]. Thus, our model could solve the strong CP problem and achieve the suitable GCU as well as realize thermal leptogenesis without gravitino problem.

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[1] M. Fukugita and T. Yanagida, Phys. Lett. B 174, 45 (1986).
[2] P. Minkowski, Phys. Lett. B 67, 421 (1977);
T. Yanagida, in Proceedings of Workshop on the Unified Theory and the Baryon Number in the Universe, Tsukuba, Japan, edited by A. Sawada and A. Sugamoto (KEK, Tsukuba, 1979), p 95;
M. Gell-Mann, P. Ramond, and R. Slansky, in Supergravity, Proceedings of Workshop, Stony Brook, New York, 1979, edited by P. Van Nieuwenhuizen and D. Z. Freedman (North-Holland, Amsterdam, 1979), p 315;
R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).
[3] For a review, see e.g. W. Buchmuller, P. Di Bari and M. Plumacher, Annals Phys. 315, 305 (2005).
[4] A. Strumia and F. Vissani, [arXiv:hep-ph/0606054].
[5] M. Maltoni, T. Schwetz, M. A. Tortola and J. W. F. Valle, New J. Phys. 6, 122 (2004);
G. L. Fogli, E. Lisi, A. Marrone and A. Palazzo, Prog. Part. Nucl. Phys. 57, 742 (2006).
[6] M. Y. Khlopov and A. D. Linde, Phys. Lett. B 138, 265 (1984);
J. R. Ellis, J. E. Kim and D. V. Nanopoulos, Phys. Lett. B 145, 181 (1984).
[7] For a recent analysis, see e.g.,
M. Kawasaki, K. Kohri, T. Moroi and A. Yotsuyanagi, Phys. Rev. D 78, 065011 (2008);
R. H. Cyburt, J. Ellis, B. D. Fields, F. Luo, K. A. Olive and V. C. Spanos, JCAP 0910, 021 (2009).
[8] E. Ma, Phys. Rev. Lett. 86, 2502 (2001).
[9] S. Gabriel and S. Nandi, Phys. Lett. B 655, 141 (2007).
[10] E. Ma, Phys. Rev. D 73, 077301 (2006).
[11] S. M. Davidson and H. E. Logan, Phys. Rev. D 80, 095008 (2009).
[12] H. E. Logan and D. MacLennan, Phys. Rev. D 81, 075016 (2010).
[13] N. Haba and M. Hirotsu, Eur. Phys. J. C 69, 481 (2010).
[14] N. Haba and O. Seto, Prog. Theor. Phys. 125, 1155 (2011).
[15] N. Haba and K. Tsumura, [arXiv:1105.1409 [hep-ph]].
[16] See, for example, Y. Ashie et al. [Super-Kamiokande Collaboration], Phys. Rev. D 71, 112005 (2005).
[17] V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. B 155, 36 (1985).
[18] L. Covi, E. Roulet and F. Vissani, Phys. Lett. B 384, 169 (1996).
[19] E. Komatsu et al. [WMAP Collaboration], Astrophys. J. Suppl. 192, 18 (2011).
[20] W. Buchmuller, P. Di Bari and M. Plumacher, Nucl Phys B 643, 367 (2002).
[21] S. Davidson and A. Ibarra, Phys. Lett. B 535, 25 (2002).
[22] M. Plumacher, Nucl. Phys. B530, 207-246 (1998).
[23] K. Nakamura et al. (Particle Data Group), J. Phys. G 37, 075021 (2010).
[24] A. Nelson, Phys. Lett. B 136, 387 (1984);
S. M. Barr, Phys. Rev. Lett. 53, 320 (1984).
[25] M. Dine, R. G. Leigh and A. Kagan, Phys. Rev. D 48, 2214 (1993).