Published by Institute of Physics Publishing for SISSA/ISAS

Received: *July 4, 2002*
Revised: *September 27, 2002*
Accepted: *October 3, 2002*

D-term inflation and neutrino mass

Daijiro Suematsu

Institute for Theoretical Physics, Kanazawa University
Kanazawa 920-1192, Japan
E-mail: suematsu@hep.s.kanazawa-u.ac.jp

Abstract: We study a D-term inflation scenario in a model extended from the minimal supersymmetric standard model (MSSM) by two additional abelian factor groups focussing on its particle physics aspects. Condensates of the fields related to the inflation can naturally give a possible solution to both the μ-problem in the MSSM and the neutrino mass through their nonrenormalizable couplings to the MSSM fields. Mixings between neutrinos and neutralinos are also induced by some of these condensates. Small neutrino masses are generated by a weak scale seesaw mechanism as a result of these mixings. Moreover, the decay of the condensates may be able to cause the leptogenesis. Usually known discrepancy between both values of a Fayet-Iliopoulos D-term which are predicted by the COBE normalization and also by an anomalous U(1) in the weakly-coupled superstring might be reconciled.

Keywords: Supersymmetry Phenomenology, Baryogenesis, Physics of the Early Universe, Neutrino Physics.

JHEP10(2002)014

Contents

1. Introduction

Recent observations in both astrophysics and particle physics provide crucial informations to astro-particle physics. Observation of the cosmic microwave background radiation (CMB) gives a strong support for the inflation [1]. Moreover, precise estimations of the spectral index of the density perturbations based on the recent observations strictly restrict the form of the inflaton potential. Particle physics models seem to be required to have a structure to induce the inflation consistent with these observations in it. On the other hand, observations of the atmospheric and solar neutrinos at Super-Kamiokande suggest the existence of nonzero neutrino masses and this fact requires some modification of the standard model (SM) [2]. From the theoretical point of view of particle physics, moreover, supersymmetrization of the SM is considered to be the most promising candidate to solve the gauge hierarchy problem, although we have no direct evidence for it still now. Thus at the present stage it seems to be natural to consider the inflation based on a suitable supersymmetric model which is motivated from a view point of particle physics such as the explanation of, for example, the neutrino mass, the μ-problem, the strong CP problem and so on [3]–[11].

Various ideas for the inflation [12]–[15] have been extended to the supersymmetric models by now. They may be classified into two types of models, that is, an F-term inflation model and a D-term inflation model. If we consider the F-term inflation, we are always suffered from the η-problem which is caused by the supergravity correction to an inflaton mass as a result of the supersymmetry breaking during the inflation [16]. As far as we do not introduce a special Kähler potential [16], there is a large correction to the inflaton mass of the order of a Hubble constant and then the inflation cannot occur. A D-term inflation [17] has been proposed as a model which can escape this η-problem. In the D-term inflation model there appear several SM singlet fields intimately related to the inflation. Although they play a crucial role in the inflation phenomena, they usually have

JHEP10(2002)014

no role in particle physics. If we can relate their existence to the important problems in particle physics, for example, the neutrino mass, the μ-problem and so on, the D-term inflation scenario might be much more promising.

In this paper we investigate some particle physics aspects of the SM singlet fields involved in the inflation in a certain D-term inflation model. The model is defined as an extension from the MSSM by two abelian factor groups $\mathrm{U}(1)_X \times \mathrm{U}(1)_A$, which are extended by an additional abelian factor group in comparison with the ordinary D-term inflation scenario. We show that the condensates of these singlet fields can be related to small neutrino masses and the μ-problem in this model. Their decay also may induce the leptogenesis. Moreover, the introduction of this new factor group might be able to relax the discrepancy between the scales of the D-term predicted by the cosmic background explorer (COBE) normalization and by the weakly-coupled superstring. The paper is organized as follows. In the next section we discuss various aspects of the inflation in this model. In section 3 we study the features as a particle physics model, especially, the origin of the neutrino masses and the μ-term. The possibility of the leptogenesis is also briefly discussed. The last section is devoted to the summary.

2. A model of D-term inflation

In our model field contents and a gauge structure are motivated by the weakly-coupled superstring models based on, for example, E_6. Such kind of the field contents and the gauge structure have been suggested to appear often in their effective theories [18]. $\mathrm{U}(1)_A$ is assumed to have a Fayet-Iliopoulos (F-I) D-term, which might be considered to be an anomalous U(1). As fields related to the inflation, we introduce a gauge singlet chiral superfield ϕ and two pairs of chiral superfields $(N, \bar{N})$ and $(S, \bar{S})$, which are the SM singlet fields but have charges of $\mathrm{U}(1)_X \times \mathrm{U}(1)_A$. A renormalizable superpotential constituted by these fields is assumed to be

$$W_1 = k\phi N\bar{N}\,, \tag{2.1}$$

where k is considered to be real and positive by a suitable redefinition of fields. A value of k is constrained on the basis of some observational requirements such as the CMB data and so on. The potential for the scalar components of these fields can be written as[1]

$$\begin{aligned} V = {} & \frac{g_X^2}{2}\left[Q_X^N\left(|\tilde{N}|^2 - |\tilde{\bar{N}}|^2\right) + Q_X^S\left(|\tilde{S}|^2 - |\tilde{\bar{S}}|^2\right)\right]^2 + \\ & + \frac{g_A^2}{2}\left[Q_A^N\left(|\tilde{N}|^2 - |\tilde{\bar{N}}|^2\right) + Q_A^S\left(|\tilde{S}|^2 - |\tilde{\bar{S}}|^2\right) + \xi_A\right]^2 + \\ & + k^2\left[|\tilde{\phi}|^2\left(|\tilde{N}|^2 + |\tilde{\bar{N}}|^2\right) + |\tilde{N}\tilde{\bar{N}}|^2\right]. \end{aligned} \tag{2.2}$$

The first two lines are D-term contributions and the last line is an F-term contribution coming from eq. (2.1). The role of $(N, \bar{N})$ and $(S, \bar{S})$ in particle physics is determined through the couplings with the MSSM matter fields. These couplings can be neglected in

[1] In the following discussion we denote a scalar component of the superfield by adding a tilde on the same character as the superfield and a spinor component by the same notation as them.

eq. (2.2) since the field values of the MSSM matter fields are fixed to be zero by the steep potential. We will come back to the effects of these couplings later and firstly we discuss features of the inflation in this model.

As similar to the usual hybrid inflation models, several fields concern the inflation in the present model.[2] A potential minimum shifts from the false vacuum to the true one around a critical value $\tilde\phi = \tilde\phi_c$ when $\tilde\phi$ varies its value. If $\tilde\phi$ initially takes a large value such as $|\tilde\phi| \gg |\tilde\phi_c|$ which is naturally expected because of the flatness of its potential, a minimum is realized at

$$\tilde N = \tilde{\bar N} = 0\,, \qquad \delta\tilde S_i^2 \equiv |\tilde S|^2 - |\tilde{\bar S}|^2 = -\frac{g_A^2 Q_A^S \xi_A}{g_X^2 Q_X^{S2} + g_A^2 Q_A^{S2}}\,, \tag{2.3}$$

where we should note that each value of $|\tilde S|$ and $|\tilde{\bar S}|$ cannot be determined here. At this minimum the potential has a constant value

$$V = \frac{1}{2} g_A^2 G \xi_A^2, \qquad G = \frac{g_X^2 Q_X^{S2}}{g_X^2 Q_X^{S2} + g_A^2 Q_A^{S2}}\,. \tag{2.4}$$

Tree level mass of $\tilde\phi$ is zero and the potential is flat in that direction. However, one-loop effect due to the interaction of eq. (2.1) induces an effective mass in the same way as the models in [17] since supersymmetry is broken by the vacuum energy at that time. In fact, although spinor components of N and $\bar N$ have the same mass $m_{N,\bar N} = k|\tilde\phi|$, the masses of the scalar components $\tilde N$ and $\tilde{\bar N}$ change differently their values depending on the value of $\tilde\phi$ as

$$\begin{aligned} m_{\tilde N}^2 &= k^2|\tilde\phi|^2 + \left[g_A^2 Q_A^N \xi_A + (g_X^2 Q_X^S Q_X^N + g_A^2 Q_A^S Q_A^N)\delta\tilde S^2\right], \\ m_{\tilde{\bar N}}^2 &= k^2|\tilde\phi|^2 - \left[g_A^2 Q_A^N \xi_A + (g_X^2 Q_X^S Q_X^N + g_A^2 Q_A^S Q_A^N)\delta\tilde S^2\right]. \end{aligned} \tag{2.5}$$

Using these relations, the effective potential [19] can be written as

$$V(\tilde\phi) = \frac{1}{2} g_A^2 G \xi_A^2 + \frac{k^4 |\tilde\phi_c|^4}{16\pi^2} \ln\frac{k^2|\tilde\phi|^2}{\Lambda^2}\,. \tag{2.6}$$

This potential makes $\tilde\phi$ slowly roll to the critical value $\tilde\phi_c$ from its initial one ($\gg |\tilde\phi_c|$). It results in a D-term inflation due to the vacuum energy (2.4) during that slow-roll period.[3] If charges of the fields are suitably selected, either $m_{\tilde N}^2$ or $m_{\tilde{\bar N}}^2$ in eq. (2.5) changes its sign at $\tilde\phi = \tilde\phi_c$ and the global minimum appears in a different point from the one of eq. (2.3).

When $|\tilde\phi| < |\tilde\phi_c|$, we find that $m_{\tilde N}^2 m_{\tilde{\bar N}}^2 < 0$ is satisfied so that either $\tilde N$ or $\tilde{\bar N}$ can have a nonzero vacuum expectation value. If we assume $m_{\tilde{\bar N}}^2 < 0$, the potential minimum shifts

[2] Although we could consider the model as an chaotic inflation model along the D-flat direction to which several fields are related such as [5, 6], it seems to be better to set it up as an hybrid inflation to avoid extremely small couplings, that is, from the view point of naturalness.

[3] If there are contributions to the above mentioned D-terms from other fields, a vacuum shift could occur to restore the supersymmetry and the inflation cannot be induced in that case [20]. However, the suitable charge assignment for such fields always make it possible to avoid such a situation. In this paper we consider such a case where their contribution to eq. (2.2) can be neglected.

JHEP10(2002)014

from a point described by eq. (2.3) to a point with $|\tilde{\bar N}| \neq 0$ and $|\tilde N| = 0$. The inflaton $\tilde\phi$ gets a large mass $m^2_{\tilde\phi} = k^2|\tilde{\bar N}|^2$ so that the slow-roll of $\tilde\phi$ stops and the inflation is expected to end at $|\tilde\phi| \sim |\tilde\phi_c|$. The critical value $\tilde\phi_c$ is estimated as

$$|\tilde\phi_c|^2 = \frac{g_A^2}{k^2}\xi_A G F, \qquad F = \frac{Q_A^N Q_X^S - Q_A^S Q_X^N}{Q_X^S}\,, \tag{2.7}$$

where in this derivation we used a value for $\delta\tilde S_i^2$ in eq. (2.3). If we use this formula in eq. (2.6), we can simplify the expression for $V(\tilde\phi)$ as

$$V(\tilde\phi) = \frac{1}{2}g_A^2 G\xi_A^2\left(1 + \frac{g_A^2}{16\pi^2}GF^2\ln\frac{k^2|\phi|^2}{\Lambda^2}\right). \tag{2.8}$$

We find that this is a simple modification by factors G and F from the inflaton potential given in [17]. The change is caused by the extension of the gauge structure and the introduction of the new SM singlet fields. On the other hand, the vacuum shift also occurs in the sector of $(\tilde S, \tilde{\bar S})$ from a state with the value of $\delta\tilde S_i^2$ in eq. (2.3) to the global minimum. The global minimum is realized at

$$\langle\tilde\phi\rangle = \langle\tilde N\rangle = 0\,, \qquad |\langle\tilde{\bar N}\rangle|^2 = \frac{\xi_A}{F}\,, \qquad \delta\tilde S_f^2 \equiv |\langle\tilde S\rangle|^2 - |\langle\tilde{\bar S}\rangle|^2 = \frac{Q_X^N}{Q_X^S}\frac{\xi_A}{F}\,. \tag{2.9}$$

At this global minimum, $V = 0$ and the supersymmetry is restored. It may be useful to note that the deviations of $\tilde{\bar N}$ and $\delta\tilde S^2$ from their true vacuum values could play a role like an Affleck-Dine (AD) condensates [21] and be relevant to the leptogenesis depending on the coupling of these fields with the MSSM matter fields. We will come back to this point later.

The superfields S, $\bar S$, N and $\bar N$ can get masses represented by

$$m^2_{\tilde N} = k^2|\langle\tilde{\bar N}\rangle|^2\,, \qquad \mathcal{M}^2 = 2\begin{pmatrix} A|\langle\tilde S\rangle|^2 & -A|\langle\tilde{\bar S}\rangle\langle\tilde S\rangle| & -C|\langle\tilde S\rangle\langle\tilde{\bar N}\rangle| \\ -A|\langle\tilde{\bar S}\rangle\langle\tilde S\rangle| & A|\langle\tilde{\bar S}\rangle|^2 & C|\langle\tilde{\bar S}\rangle\langle\tilde{\bar N}\rangle| \\ -C|\langle\tilde S\rangle\langle\tilde{\bar N}\rangle| & C|\langle\tilde{\bar S}\rangle\langle\tilde{\bar N}\rangle| & B|\langle\tilde{\bar N}\rangle|^2 \end{pmatrix}, \tag{2.10}$$

where A, B and C are defined as

$$\begin{aligned} A &= (g_X^2 Q_X^{S2} + g_A^2 Q_A^{S2})\,, \qquad B = (g_X^2 Q_X^{N2} + g_A^2 Q_A^{N2})\,, \\ C &= (g_X^2 Q_X^S Q_X^N + g_A^2 Q_A^S Q_A^N)\,. \end{aligned} \tag{2.11}$$

The matrix $\mathcal{M}^2$ is written by the basis of $(S,\ \bar S,\ \bar N)$. The masses of spinor components of $(S,\ \bar S,\ \bar N)$ are generated through the mixings with gauginos of $\mathrm{U}(1)_X$ and $\mathrm{U}(1)_A$. Since $m^2_{\tilde{\bar N}}$ in eq. (2.10) is sufficiently large, the value of $\langle\tilde{\bar N}\rangle$ is kept in the same one even after the introduction of the soft supersymmetry breaking effects. On the other hand, mass eigenstates of $\mathcal{M}^2$ have an interesting feature. One of them is massless and it is composed of S and $\bar S$ such as $\Phi_0 \propto \langle\tilde{\bar S}\rangle S + \langle\tilde S\rangle\bar S$. The fact that each value of $\langle\tilde S\rangle$ and $\langle\tilde{\bar S}\rangle$ is not determined in eq. (2.9) comes from the existence of this mode. The other mass eigenstates Φ_1 and Φ_2 have masses of $O(\sqrt{\xi_A/F})$ and are composed as the linear combinations of $\langle\tilde{\bar S}\rangle S - \langle\tilde S\rangle\bar S$ and $\bar N$. A mass of the massless state Φ_0 is considered to be supplied through the soft supersymmetry breaking effects and also the radiative correction on it, which are considered to be of the order of the weak scale. When its mass is generated, the scalar component of $\tilde{\Phi}_0$ starts to oscillate around the true vacuum keeping the condition for $\delta\tilde{S}_f^2$ in eq. (2.9) and its decay is expected to induce the reheating. Since the mass of Φ_0 is of the order of the weak scale, it is also expected to play a certain role in the particle physics phenomenology at a low energy scale.

By now we have briefly sketched the inflation scenario in this model. We discuss more detailed quantitative features here. The slow-roll parameters of the inflation defined as $\epsilon = (1/2)\tilde{M}_{\rm pl}^2(V'/V)^2$ and $\eta = \tilde{M}_{\rm pl}^2 V''/V$ should satisfy the conditions $\epsilon,\ |\eta| \lesssim 1$ during the slow-roll period.[4] The values of these slow-roll parameters at $|\tilde{\phi}| = |\tilde{\phi}_c|$ is estimated by using $V(\tilde{\phi})$ in eq. (2.8) as

$$\epsilon = \frac{k^2 g_A^2}{128\pi^4}\frac{\tilde{M}_{\rm pl}^2}{\xi_A}GF^3\,, \qquad \eta = -\frac{k^2}{8\pi^2}\frac{\tilde{M}_{\rm pl}^2}{\xi_A}F\,. \tag{2.12}$$

The end of the slow-roll inflation is determined as the earlier period of the realization of $|\tilde{\phi}| \simeq |\tilde{\phi}_c|$ or $|\eta| \simeq 1$. If we write the value of $\tilde{\phi}$ at that time as $\tilde{\phi}_{\rm end}$, we find that $|\tilde{\phi}_c| \simeq |\tilde{\phi}_{\rm end}|$ is realized when $k^2F \simeq 8\pi^2\xi_A/\tilde{M}_{\rm pl}^2$ is satisfied. The e-folds during the slow-roll inflation is written by using the slow-roll approximation as

$$N(\tilde{\phi}) = \int_{\tilde{\phi}_{\rm end}}^{\tilde{\phi}} \frac{1}{\tilde{M}_{\rm pl}^2}\frac{V}{V'}\,d\tilde{\phi} = \frac{4\pi^2}{g_A^2}\frac{1}{GF^2}\frac{1}{\tilde{M}_{\rm pl}^2}\left(|\tilde{\phi}|^2 - |\tilde{\phi}_{\rm end}|^2\right), \tag{2.13}$$

and then $|\tilde{\phi}|^2 \simeq (1/4\pi^2)N(\tilde{\phi})g_A^2GF^2\tilde{M}_{\rm pl}^2$ where we used the relation $|\tilde{\phi}| \gg |\tilde{\phi}_{\rm end}|$. In eq. (2.13), $\tilde{\phi}$ is a value at the period when a relevant cosmological scale of the universe leaves the horizon.[5]

The CMB data requires that ξ_A should satisfy a suitable condition. The COBE normalization for the density perturbations is given as $\delta_H = \sqrt{V/(150\pi^2\tilde{M}_{\rm pl}^4\epsilon)} \sim 1.95\times 10^{-5}$ [1], where V and V' are also estimated at the epoch where the COBE scale leaves the horizon. By using eq. (2.13), we find that this imposes the condition[6]

$$\xi_A \simeq 8.5\times 10^{-6}\left(\frac{50}{N(\tilde{\phi})}\right)^{1/2} F\tilde{M}_{\rm pl}^2\,. \tag{2.14}$$

Combining a result given below eq. (2.12) and eq. (2.14) which is the result of the COBE normalization, we find that $|\tilde{\phi}_c| \simeq |\tilde{\phi}_{\rm end}|$ is realized for k such as $k_0 \equiv 0.025\,(50/N)^{1/2}$. If k is smaller than this value, $|\tilde{\phi}_c| > |\tilde{\phi}_{\rm end}|$ is satisfied and the fast-roll inflation period appears additionally after the end of the slow-roll period. To make the situation simple, we assume $k \simeq k_0$ in the following part as far as we do not mention it.

---

[4] $\tilde{M}_{\rm pl}$ is a reduced Planck mass defined as $\tilde{M}_{\rm pl} = M_{\rm pl}/\sqrt{8\pi}$.

[5] Unless $g_A^2GF^2$ is sufficiently smaller than 1, we cannot have sufficient e-folds such as $N \sim 50$ keeping a condition $|\tilde{\phi}| \ll \tilde{M}_{\rm pl}$, which guarantees the validity to treat the system by the field theory. However, as we will see next, it cannot be small enough.

[6] We ignore any gravitational wave contributions since it can be safely neglected in the present model.

JHEP10(2002)014

As the usual hybrid inflation and D-term inflation, this model predicts an almost flat spectrum of the density perturbations. In fact, by using eqs. (2.13) and (2.14) the spectral index is estimated as [26]

$$n - 1 = 2\eta - 6\epsilon \simeq -\frac{1}{N(\tilde{\phi})}\left(1 + \frac{3}{16\pi^2}g_A^2 G F^2\right). \tag{2.15}$$

This shows a good agreement with the analysis of the recent observations [25]. Although the second term in the parentheses of eq. (2.15) might be rather large and produce a non-negligible effect depending on the value of F, it is natural to consider that $g_A^2 G F^2$ is at most $O(1)$. Then this term is expected to be small enough. The deviation of the spectral index n from 1 is estimated as $-1/N(\tilde{\phi})$, which is the well-known feature of the D-term inflation [26]. This feature may be useful to discriminate this model from others also.

If $F = O(1)$ is assumed to be satisfied, for example, eq. (2.14) shows that ξ_A takes a value of order 10^{16} GeV. On the other hand, if the F-I D-term ξ_A is assumed to have its origin in the anomalous U(1) in the weakly-coupled superstring, ξ_A can be written as [22]

$$\xi_A = \frac{\mathrm{Tr}\ Q_A}{192\pi^2} g_{\mathrm{st}}^2 \tilde{M}_{\mathrm{pl}}^2 . \tag{2.16}$$

It takes a value in the range 10^{17-18} GeV for $g_{\mathrm{st}} \sim 0.1 - 1$ and Tr $Q_A \sim O(100)$, which are usually considered to be the typical values. In the usual D-term inflation model, there is a discrepancy between this and the above mentioned value required by the CMB data [23]. In this model, however, the effective ξ_A which dominates the vacuum energy contributing to the density perturbations has an additional factor F as seen in eq. (2.14). Because of this factor, as suggested in [20], there may be a potential possibility to reconcile this discrepancy[7] as far as F takes a value $F \sim 6 \times 10^3 (N(\tilde{\phi})/50)^{1/2} g_{\mathrm{st}}^2$. A rather large F seems to be obtainable if $Q_A^S \gg Q_X^S$ and $g_X \simeq g_A(= g_{\mathrm{st}})$ is satisfied. In such a case, both $G \ll 1$ and $F \gg 1$ can be satisfied as far as $Q_X^N > O(1)$, as can be easily seen from their expressions in eqs. (2.4) and (2.7). Under this situation, $g_A^2 G F^2$ can be estimated as $\sim g_X^2 Q_X^{N2}$ and it cannot take a small value so as to make $|\tilde{\phi}|$ sufficiently small in comparison with $\tilde{M}_{\mathrm{pl}}$. We must consider the case with $|\tilde{\phi}| \lesssim \tilde{M}_{\mathrm{pl}}$ and the validity to treat the system by the field theory seems to be marginal. It depends on the value of g_{st}. If the Yukawa coupling k in W_1 is smaller than k_0, however, there is an additional fast-roll inflation after this slow-roll inflation. In this case $N(\tilde{\phi})$ can be smaller and as a result F can be somewhat smaller. Thus the smaller values of k and g_{st} seems to give a promising possibility for the reconciliation or the relaxation of the discrepancy between the required ξ_A values . In such a situation, $|\tilde{\phi}| < \tilde{M}_{\mathrm{pl}}$ is satisfied and $|n-1|$ becomes somewhat larger because of the smaller value of $N(\tilde{\phi})$ required for the slow-roll inflation.

Next we consider the reheating in this model. As mentioned in the previous part, $\tilde{\phi}$ and the scalar components of two massive states Φ_1, Φ_2 start to oscillate around the global minimum when $|\tilde{\phi}| \lesssim |\tilde{\phi}_c|$ is realized. Preheating and reheating processes are expected to

[7] We do not consider a kinetic term mixing between $\mathrm{U}(1)_X$ and $\mathrm{U}(1)_A$ [24]. If we take it into account, the situation for this discrepancy seems to be worse.

JHEP10(2002)014

proceed through these oscillations. On the preheating related to $\tilde{\phi}$ and $\tilde{\bar{N}}$ in this model, the study in [27] is very useful. The hybrid inflation model with the scalar potential

$$V(\phi,\sigma)=\frac{1}{4\lambda}(\lambda\sigma^2-M^2)^2+\frac{1}{2}g^2\phi^2\sigma^2+\frac{1}{2}m^2\phi^2 \tag{2.17}$$

has a very similar structure to the scalar sector in the present model. If we use the following replacement in the scalar potential in [27]:

$$\phi\to\tilde{\phi}\,,\quad \sigma\to\tilde{\bar{N}}\,,\quad M^2\to -g_A^2Q_A^N\xi_A\,,\quad g^2\to k^2\,,\quad \lambda\to g_X^2Q_X^{N2}+g_A^2Q_A^{N2}\,, \tag{2.18}$$

a part of the scalar potential (2.2) can be obtained. We find that our model corresponds to a case of $\lambda\gg g^2$ in [27] for this sector by taking account of the previous discussion on the couplings. Following their study, the oscillation energy is expected to be concentrated immediately to the oscillation of ϕ in that case. They also shows that there is no effective preheating and no significant particle production of ϕ and σ through its oscillation. We expect that their results can be applied to our model and the preheating effect can be also neglected here. Then the deviation of $\tilde{\Phi}_1$ and $\tilde{\Phi}_2$ which are composed of $\tilde{S}$, $\tilde{\bar{S}}$ and $\tilde{\bar{N}}$ from their true vacuum values during the inflation cannot cause the oscillation since their oscillation energies are expected to be instantaneously transferred into the oscillation of $\tilde{\phi}$ through the coupling with $\tilde{\bar{N}}$. We need to consider only the reheating due to the usual perturbative $\tilde{\phi}$ decay. The decays of the inflaton $\tilde{\phi}$ into N and $\Phi_{1,2}$ is expected to be kinematically prohibited in the present model since $m^2_{\Phi_1}$, $m^2_{\Phi_2}>m^2_{\tilde{\phi}}$ is satisfied following the discussions below eqs. (2.10) and (2.16). Moreover, since Φ_0 has no $\bar{N}$ component in it, $\tilde{\phi}$ cannot decay into Φ_0 through the coupling $k\phi\bar{N}N$. Thus we cannot expect the reheating by the $\tilde{\phi}$-decay into the light MSSM particles through the couplings in the renormalizable superpotential. However, in general, its oscillation can decay into the MSSM light fields through the allowed nonrenormalizable couplings contained in the superpotential which may be expressed as

$$W_2=c\,\frac{\phi\psi^r}{\tilde{M}_{\rm pl}^{r-2}}\,, \tag{2.19}$$

where ψ stands for the MSSM chiral superfields and c is assumed to be $O(1)$. If $\tilde{\phi}$ has a decay width $\Gamma_{\tilde{\phi}}$ through this type of coupling, the reheating temperature can be estimated by assuming the Hubble parameter H to be $H\sim\Gamma_{\tilde{\phi}}$ as [32]

$$T_{\rm RH}^{\tilde{\phi}}\sim\frac{c\;k^{r-(3/2)}}{g_*^{1/4}(3.6\times10^2)^{r-4}}\left(\frac{P_r}{P_4}\right)^{1/2}\times1.3\times10^{10}\,{\rm GeV}\,, \tag{2.20}$$

where g_* is an effective number of the relativistic degrees of freedom at this period and P_r/P_4 represents the ratio of the phase space factors between r- and 4-body decays, which decreases with the increase of r. As far as $r\geq4$ is satisfied, the constraint coming from the gravitino problem can be satisfied. The decay of $\tilde{\phi}$ occurs at $H\lesssim M_{\rm susy}$ depending on the value of r (≥4), where $M_{\rm susy}$ stands for a typical soft supersymmetry breaking scale of order 1 TeV. The value of r and the couplings of S and $\bar{S}$ to the light MSSM fields determine whether the condensate of $\tilde{\phi}$ can decay earlier than the deviation of Φ_0

from the true vacuum or not. This largely affects the baryogenesis scenario in this model. For this study we need to fix the couplings of S and $\bar{S}$ with the MSSM fields in the nonrenormalizable terms of the superpotential.

3. Implications for particle physics

We need to introduce additional interactions of the SM singlet chiral superfields (N, $\bar{N}$) and (S, $\bar{S}$) with the MSSM chiral superfields in the superpotential without modifying the features of the inflation in order to fix their role in the particle physics phenomenology. Taking account that both $\tilde{S}$ and $\tilde{\bar{S}}$ can have nonzero vacuum expectation values (VEVs) at the true vacuum as shown in eq. (2.9), we consider the following terms as the dominant ones in the superpotential,

$$W_3 = h\frac{(S\bar{S})^\ell}{\tilde{M}_{\rm pl}^{2\ell}}L_\alpha H_2\bar{N} + \lambda\frac{(S\bar{S})^n}{\tilde{M}_{\rm pl}^{2n}}SH_1H_2 + \gamma\frac{(S\bar{S})^m}{\tilde{M}_{\rm pl}^{2m-3}}\,, \tag{3.1}$$

where the powers ℓ, m and n are assumed to satisfy $\ell,\ n \geq 0$ and $m \geq 2$, respectively. The coupling constants h, λ and γ should be assumed to be $O(1)$ at least when the couplings are nonrenormalizable.[8] A chiral superfield L_α stands for a doublet lepton and $H_1,\ H_2$ are the usual doublet Higgs chiral superfields. The addition of these terms does not affect the above discussed inflation scenario. It is useful to note that the usual D-term flatness along the well-known directions H_1H_2 and LH_2 in the MSSM is lost because of the additional U(1) interactions. Thus, differently from the MSSM, the Affleck-Dine leptogenesis [33] along the latter direction cannot be expected in this model. The model is naturally expected to have the soft supersymmetry breaking parameters corresponding to the superpotential W_3 such as

$$\mathcal{L}_{\rm soft} = A_h\frac{(\tilde{S}\tilde{\bar{S}})^\ell}{\tilde{M}_{\rm pl}^{2\ell}}\tilde{L}_\alpha\tilde{H}_2\tilde{\bar{N}} + A_\lambda\frac{(\tilde{S}\tilde{\bar{S}})^n}{\tilde{M}_{\rm pl}^{2n}}\tilde{S}\tilde{H}_1\tilde{H}_2 + \cdots\,, \tag{3.2}$$

where A_h and A_λ are assumed to be $O(M_{\rm susy})$ and complex.

3.1 μ-term and neutrino masses

The structure of the superpotential W_3 crucially affects the phenomenology at a low energy region. For example, the second term in W_3 may be considered to induce an effective μ-term of the MSSM in the form as $\mu = (\lambda/\tilde{M}_{\rm pl}^{2n})\langle\tilde{S}\rangle^{n+1}\langle\tilde{\bar{S}}\rangle^n$ if both $\tilde{S}$ and $\tilde{\bar{S}}$ get suitable VEVs. The ratio $\tan\alpha_S \equiv \langle\tilde{S}\rangle/\langle\tilde{\bar{S}}\rangle$ determines the composition of Φ_0 and also a value of the reheating temperature $T_{\rm RH}^{\tilde{S}}$ due to the decay of the condensate of $\tilde{\Phi}_0$ as seen later. As mentioned below eq. (2.10), there is a massless state Φ_0 which contains S with the weight $\cos\alpha_S$. The mass of $\tilde{\Phi}_0$ has the dominant contributions only from the soft supersymmetry breaking squared mass, radiative corrections based on the supersymmetry breaking effects and also the last term in W_3. If its squared mass becomes negative by the radiative

[8] We do not consider the generation dependence of h, for simplicity. The third term is introduced to stabilize the potential for $\tilde{S}$ and $\tilde{\bar{S}}$. It could contribute the inflation potential (2.2). However, we can check its effect is not sufficiently large to change the present inflation scenario.

correction through the couplings with extra colored chiral superfields g and $\bar{g}$ such as $Sg\bar{g}$ in the superpotential [28, 29] or by the suitable supersymmetry breaking [30], for example, the values of $\langle \tilde{S} \rangle$ and $\langle \tilde{\bar{S}} \rangle$ are expected to be shifted into the true vacuum ones satisfying the relation for $\delta \tilde{S}_f^2$ in eq. (2.9). Thus, as far as we do not consider the large cancellation between the contributions from $\langle \tilde{S} \rangle$ and $\langle \tilde{\bar{S}} \rangle$, either $\langle \tilde{S} \rangle$ or $\langle \tilde{\bar{S}} \rangle$ are at least required to be an almost GUT scale.[9] The value of $\tan\alpha_S$ is expected to take a rather wide range value [29] depending on the details of the model, for example, the extra matter contents coupled with S and $\bar{S}$ and the power m, and also the nature of the symmetry breaking in this sector. There seem to be two possibilities to determine the values of $\langle \tilde{S} \rangle$ and $\langle \tilde{\bar{S}} \rangle$. One is the case where the third term in W_3 plays a crucial role in their determination. The other case is based on the purely radiative effects. In that case the value of the power m is large enough and then the third term in W_3 can have only the much smaller effect than the radiative correction [29]. Although the latter possibility needs the numerical study to get some results,[10] the former one allows us to have a rough estimation of the symmetry breaking scale.

For this instruction, let us present an example which is found to be interesting from some phenomenological reasons as shown later. We consider the situation such that Φ_0 is dominated by S and then $\langle \tilde{S} \rangle \ll \langle \tilde{\bar{S}} \rangle \sim O(10^{15})$ GeV is satisfied. We also require the realization of a suitable μ value such as $O(1)$ TeV. Then, if we assume that the third term in W_3 can fix the values of $\langle \tilde{S} \rangle$ and $\langle \tilde{\bar{S}} \rangle$, we find that $\tan\alpha_S = \lambda^{-1/(n+1)}10^{6(n-2)/(n+1)}$ should be satisfied and the allowed value of n is restricted to zero or one. As seen later, the phenomenological condition for the reheating temperature can give a severe constraint on the value of n. Here let us assume the masses of $\tilde{S}$ and $\tilde{\bar{S}}$ to be $m^2_{\tilde{S}} < 0$ and $m^2_{\tilde{\bar{S}}} < 0$ due to some radiative effects. The potential minimization gives

$$\begin{aligned} |\langle \tilde{S} \rangle| &\sim \tilde{M}_{\rm pl}^{\frac{2m-3}{2(m-1)}} |m_{\tilde{S}}|^{-1/2} |m_{\tilde{\bar{S}}}|^{\frac{m}{2(m-1)}} , \\ |\langle \tilde{\bar{S}} \rangle| &\sim \tilde{M}_{\rm pl}^{\frac{2m-3}{2(m-1)}} |m_{\tilde{S}}|^{1/2} |m_{\tilde{\bar{S}}}|^{\frac{2-m}{2(m-1)}} . \end{aligned} \qquad (3.3)$$

From this result, we find $\tan\alpha_S \sim |m_{\tilde{\bar{S}}}|/|m_{\tilde{S}}|$. The situation with the suitable value of $\tan\alpha_S$ is realized in the case of $m = 3$ for the appropriately tuned scalar masses in the case of $n = 1$.[11] The situation may be somewhat different from this in the pure radiative symmetry breaking case.

In the model defined by $W_1 + W_2 + W_3$ the singlet field $\bar{N}$ can be understood as a charge conjugated field of a right-handed neutrino and it has a lepton number -1. Thus we can consider a neutrino mass generation based on the first coupling in W_3. Since we have a large Majorana mass for the right-handed neutrino $\bar{N}$ as shown in eq. (2.10), the ordinary seesaw mechanism [34] might be expected to work. In fact, $\bar{N}$ gets a large Majorana mass of $O(|\langle \tilde{\bar{N}} \rangle|)$ and then we can obtain a small neutrino mass, which is estimated for the $\ell = 0$

[9]This symmetry breaking scale is somewhat larger than the one assumed in the similar model where the μ-problem and the strong CP problem have been considered [6]–[9].

[10]This study is beyond the scope of the paper and we do not discuss this case further.

[11]In the case of $n = 0$, we have $\tan\alpha_S = 10^{-12}$ and its realization requires extremely tuned $m^2_{\tilde{S}}$ and $m^2_{\tilde{\bar{S}}}$.

case, as

$$m_\nu \sim \frac{h^2 |\langle \tilde{H}_2 \rangle|^2}{|\langle \tilde{\tilde{N}} \rangle|} . \tag{3.4}$$

If we impose $m_\nu \sim O(10^{-1})\,\mathrm{eV}$ based on the atmospheric neutrino analysis, we find that h should satisfy $h \sim 0.01$ by using eq. (2.9). However, the same term in eq. (3.1) spontaneously induces a bilinear R-parity violating term $\varepsilon L H_2$ with $\varepsilon = h\langle \tilde{\tilde{N}} \rangle$. If we take $h \sim 0.01$, ε becomes too large and such a large R-parity violating term is forbidden by various phenomenological constraints [35]. This means that the magnitude of the Yukawa coupling constant h required by the phenomenological condition coming from the R-parity violation largely contradicts with the one required by the neutrino mass.

Fortunately, we can consider another possibility for the generation of neutrino masses. Neutrino masses may be produced by neutrino-neutralino mixings as discussed in [36]–[38]. If ε is sufficiently small as a result of the higher orderness of the relevant term in the superpotential W_3, a weak scale seesaw mechanism is expected to induce neutrino masses as

$$m_\nu \simeq \frac{\varepsilon^2}{\mu} \simeq (\tan \alpha_S)^{2\ell-n-1} \left(\frac{|\langle \tilde{\tilde{S}} \rangle|}{\tilde{M}_{\rm pl}} \right)^{4\ell-2n} \frac{|\langle \tilde{\tilde{N}} \rangle|^2}{|\langle \tilde{\tilde{S}} \rangle|} , \tag{3.5}$$

where we use that h and λ are $O(1)$. Since $|\langle \tilde{\tilde{S}} \rangle| \sim \langle |\tilde{\tilde{N}} \rangle| \sim O(10^{15})\,\mathrm{GeV}$ should be satisfied in this inflation model, we find that m_ν can be in the suitable range of $O(10^{-1})$ eV for the atmospheric neutrino observation by taking $\ell = 2$ and $n = 1$. In this case $|\varepsilon| \sim 10^{-6}|\mu|$ and $\mu = O(1)$ TeV are satisfied.[12] Moreover, this ratio of ε and μ is independent of the value of $\langle \tilde{S} \rangle$ in this case, which is a useful fact for the later discussion. Since the extreme smallness of the effective Yukawa coupling of the neutrino due to the nonrenormalizablity makes the above mentioned ordinary seesaw mass (3.4) negligible, the neutrino mass and mixing should be considered on the basis of the neutrino-neutralino mixing along the discussions in [36]–[38]. It is interesting that in this scenario the $\bar{N}$ sector is not necessarily required to be extended into the multi-generation in order to explain the required pattern of masses and mixings of neutrinos. Thus we can directly apply the present inflation scenario to the explanation of small neutrino masses.

In this case it is expected the appearance of a physical Nambu-Goldstone boson after the singlet fields $\tilde{\tilde{N}}$, $\tilde{S}$ and $\tilde{\tilde{S}}$ get the VEVs. Its physical feature is completely dependent on its composition. If it is dominantly composed of $\tilde{\tilde{N}}$, it may behave like a Majoron. However, it seems to be hard to find its physical implication since its effective couplings to the ordinary matter fields are too weak. On the other hand, if it is dominated by $\tilde{S}$ or $\tilde{\tilde{S}}$, it may behave like an axion related to the strong CP problem. In such a case its consistency with various constraints depends on the details of the model and its study seems to be beyond the scope of this work. We will discuss this point in other place.

If we exchange the role of $\tilde{N}$ and $\tilde{\tilde{N}}$ in eq. (2.9) by modifying their charge assignments for $\mathrm{U}(1)_X \times \mathrm{U}(1)_A$, there appears a new possibility for the generation of neutrino masses.

[12]Here we give a naive estimation of ε, as an example. However, as discussed in [39], the neutrino mass cannot be expressed in such a simple form as ε^2/μ in the present model. Following their detailed analysis, we can have $O(10^{-1})$ eV neutrino mass even for $|\varepsilon| \sim 10^{-4\sim-3}|\mu|$.

In this case $\tilde{\bar{N}}$ has no VEV at the global minimum and then we have no bilinear R-parity violating term. However, its fermion partner gets a large Majorana mass $k|\langle\tilde{N}\rangle|$ through the mixing with ϕ in W_1. Small neutrino masses can be explained by the ordinary seesaw mechanism in the case of $\ell = 0$ which corresponds to the renormalizable neutrino Yukawa coupling. They can be estimated as

$$m_\nu \sim \frac{h^2|\langle\tilde{H}_2\rangle|^2}{k|\langle\tilde{N}\rangle|}\,. \tag{3.6}$$

If we impose $m_\nu \sim O(10^{-1})$ eV, h should satisfy $h \sim 10^{-2}k^{1/2}$. In this case, however, unless we introduce additional $\bar{N}$'s, we need some new mechanism to explain the required pattern of masses and mixings of neutrinos. On the other hand, if we introduce new $\bar{N}$'s, some modification will be required in the above inflation scenario. Because of these reasons the neutrino mass generation based on the neutrino-neutralino mixing seems to be the most promising in the present inflation scenario.

We find that the superpotential W_3 with $\ell = 2$ and $n = 1$ is promising from the μ-problem and neutrino masses in the present scenario. This kind of favorable structure of the superpotential $W_1 + W_2 + W_3$ is reproduced consistently with the Yukawa couplings in the MSSM by introducing a suitable discrete symmetry. For example, as a simple candidate for such a discrete symmetry, we can consider the case in which the MSSM matter fields have the same charge. We take the symmetry as Z_{24} and adopt the following charge assignment:[13]

$$\begin{aligned} &Q_Z(Q_\alpha, L_\alpha, \bar{U}_\alpha, \bar{D}_\alpha, \bar{E}_\alpha) = 3\,, \qquad Q_Z(H_1, H_2) = -6\,, \\ &Q_Z(N) = 1\,, \quad Q_Z(\bar{N}) = -13\,, \quad Q_Z(S, \bar{S}) = 4\,, \quad Q_Z(\phi) = 12\,, \end{aligned} \tag{3.7}$$

where α is a generation index. This discrete symmetry realizes the superpotential W_1 and W_3 with $\ell = 2$, $n = 1$, $m = 3$ and also W_2 with $r = 4$.

3.2 Generation of baryon number asymmetry

Finally we discuss the baryogenesis in this model. After the inflation, the condensate of $\tilde{\bar{N}}$ (or $\tilde{\Phi}_{1,2}$) immediately reduces its value to a true vacuum one and it cannot effectively oscillate as mentioned before. As a result, we cannot use its decay for the leptogenesis. On the other hand, since a scalar component of the massless state Φ_0 is deviated from its true vacuum value after its squared mass becomes negative, $\tilde{\Phi}_0$ starts to oscillate around its true vacuum value at $H \sim m_{\tilde{\Phi}_0}$. The leptogenesis may be expected to occur through the decay of $\tilde{\Phi}_0$ accompanied by its oscillation. The oscillation of $\tilde{\Phi}_0$ which contains $\tilde{S}$ and $\tilde{\bar{S}}$ decays into L and H_2 or H_1 and H_2 through the first or second couplings in eq. (3.1). Following this decay, L, H_1 and H_2 are also transfered into the other light SM fields and they also contribute to the thermalization of the universe. The decay width of $\tilde{S}$ is estimated as $\Gamma_{\tilde{S}} \sim (1/4\pi)f^2 M_{\rm susy}$ because of $m_{\tilde{\Phi}_0} \sim M_{\rm susy}$. The effective coupling f for

[13] There are many other possibilities. We only present this example to show that the control of the superpotential by the discrete symmetry is not so difficult in this model.

JHEP10(2002)014

the $\tilde{S}$ decay into each mode should be understood as $h_{\rm eff} = (h/\tilde{M}_{\rm pl}^{2\ell})(\tan\alpha_S)^{\ell-1}\langle\tilde{\bar{S}}\rangle^{2\ell-1}\langle\tilde{N}\rangle$ or $\lambda_{\rm eff} = (\lambda/\tilde{M}_{\rm pl}^{2n})(\tan\alpha_S)^n\langle\tilde{\bar{S}}\rangle^{2n}$, respectively. Depending on the relative strength of these effective couplings $h_{\rm eff}$ and $\lambda_{\rm eff}$, the main decay mode of the $\tilde{\Phi}_0$ condensate is determined. Since we found that W_3 with $\ell = 2$ and $n = 1$ is favorable for the explanation of both the μ-problem and the small neutrino masses in the previous discussion, we will concentrate our attention on this possibility. In this case the $\tilde{\Phi}_0$ condensate mainly composed of $\tilde{S}$ decays into H_1 and H_2 through the second effective coupling $\lambda_{\rm eff}\tilde{S}H_1H_2$. In this decay the reheating temperature can be estimated as

$$T_{\rm RH}^{\tilde{S}} \simeq \frac{1.7}{g_*^{1/4}}\sqrt{\Gamma_{\tilde{S}}M_{\rm pl}} \sim \frac{0.48}{g_*^{1/4}}\lambda_{\rm eff}\sqrt{\tilde{M}_{\rm pl}M_{\rm susy}}\,. \tag{3.8}$$

This reheating temperature should satisfy a certain upper bound to prohibit the over production of gravitinos [31] and also it is high enough to allow us to use the electroweak sphaleron interaction for the production of the baryon number asymmetry. These requirement give a condition on $T_{\rm RH}^{\tilde{S}}$ such as 10^2 GeV$\lesssim T_{\rm RH}^{\tilde{S}} \lesssim 10^9$ GeV. This means that $\lambda_{\rm eff}$ has to satisfy $10^{-9} \lesssim \lambda_{\rm eff} \lesssim 10^{-2}$. This condition can be satisfied in the above discussed $n = 1$ case where the nonrenormalizable term determines the vacuum. In fact, since we have $|\langle\tilde{\bar{S}}\rangle| = O(10^{15})\,$GeV and $\tan\alpha_S = 10^{-3}$ in this case, $\lambda_{\rm eff} = O(10^{-9})$ and $T_{\rm RH}^{\tilde{S}} \sim O(10^2)\,$GeV is realized for $g_* = O(10^2)$. The pure radiative symmetry breaking case may be also expected to be able to satisfy this condition similarly.

We may be possible to consider several baryogenesis scenarios through the electroweak sphaleron interaction. The first possibility is the usual electroweak baryogenesis scenario in the electroweak phase transition. The μ term is effectively generated through the coupling with the SM singlet fields, the phase transition can have the stronger first orderness in comparison with the MSSM case because of the existence of the trilinear coupling of the Higgs scalar fields as in the case of the NMSSM. This possibility will, however, be expected only in the case where either $\langle\tilde{S}\rangle$ or $\langle\tilde{\bar{S}}\rangle$ takes a value not so far from the weak scale.

The second possibility may be considered on the basis of the decay of the oscillation of $\tilde{\Phi}_0$. In this case the problem is whether the decay product can bring the baryon number asymmetry or not. Since the oscillation of $\tilde{\Phi}_0$ is non-thermal, this decay proceeds in the way of the out-of-thermal equilibrium. Thus, if there are some CP violating complex phases and $\Gamma_{\tilde{S}} < H_{T=m_{\tilde{\Phi}_0}}$ is satisfied, we can expect the production of the asymmetry between the Higgsino $H_{1,2}$ number density and the anti-Higgsino $H_{1,2}^c$ number density through the decay by the second term in eq. (3.1). This out-of-equilibrium condition restricts the region of $\lambda_{\rm eff}$ obtained from the condition for $T_{\rm RH}^{\tilde{S}}$ further into

$$10^{-9} \lesssim \lambda_{\rm eff} \lesssim 10^{-7}\,. \tag{3.9}$$

We should note that our concrete example satisfies this condition. The pure radiative symmetry breaking may be able to loose this condition somewhat [29]. The Higgsino asymmetry ϵ_S can be produced dominantly by the interference term between a tree diagram of $\tilde{S} \to H_1H_2$ and a one-loop contribution coming from a diagram with an A_λ vertex and a gaugino $\lambda_{1,2}$ internal line. It can be roughly estimated as [40]

$$\epsilon_S \equiv \frac{\Gamma_u(\tilde{S} \to H_1 H_2) - \Gamma_u(\tilde{S}^c \to H_1^c H_2^c)}{\Gamma_u(\tilde{S} \to H_1 H_2) + \Gamma_u(\tilde{S}^c \to H_1^c H_2^c)} \sim \frac{g_2^2 + g_1^2}{16\pi} \frac{|A_\lambda|}{m_{\tilde{\Phi}_0}} \sin\delta \,, \tag{3.10}$$

where $\delta = \arg(A_\lambda)$. Thus, by using the relations $\rho_{\rm tot} = sT_{\rm RH}$ and $\rho_{\tilde{\Phi}_0} = n_{\tilde{\Phi}_0} m_{\tilde{\Phi}_0}$, the ratio of the total asymmetry of the H_2 number density to the entropy density can be estimated as

$$\frac{n_{\tilde{H}_2} - n^c_{\tilde{H}_2}}{s} = \epsilon_S \, \frac{\rho_{\tilde{\Phi}_0}}{\rho_{\rm tot}} \frac{T_{\rm RH}^{\tilde{\Phi}_0}}{m_{\tilde{\Phi}_0}} \simeq \frac{g_2^2 + g_1^2}{16\pi} \frac{T_{\rm RH}^{\tilde{\Phi}_0}}{M_{\rm susy}} \sin\delta \,, \tag{3.11}$$

where $\rho_{\tilde{\Phi}_0}$ and $\rho_{\rm tot}$ are the energy density stored in the $\tilde{\Phi}_0$ oscillation and the total energy density of the universe, respectively. The energy density of the universe is assumed to be occupied by the oscillation energy of $\tilde{\Phi}_0$ at this period and then $\rho_{\rm tot} \sim \rho_{\tilde{\Phi}_0}$.

Here we should remind that we have a bilinear R-parity violating term which comes from the first term in eq. (3.1). This term induces the mixing of $O(|\varepsilon|/|\mu|)$ between L_α and H_2. If the lightest neutralinos and charginos are dominantly composed of the H_2 components, the produced Higgsino asymmetry ϵ_S is expected to be transferred into the lepton number asymmetry through this lepton number violating mixing $\varepsilon L_\alpha H_2$. We consider the case where these neutralino and chargino need longer time to reach the thermal equilibrium through various interactions in comparison with the typical sphaleron interaction time at this period.[14] Then there is sufficient time for the sphaleron interaction in the thermal equilibrium to convert the Higgsino H_2 asymmetry into the baryon asymmetry effectively through this lepton number violating mixing. By using this and eq. (3.11), the baryon number asymmetry generated through this process can be expected to be

$$Y_B = \frac{n_B - n_{\bar{B}}}{s} \simeq c_s \kappa \, \frac{|\varepsilon|^2}{|\mu|^2} \left(\frac{g_2^2 + g_1^2}{16\pi} \frac{T_{\rm RH}^{\tilde{\Phi}_0}}{M_{\rm susy}} \sin\delta \right) \sim 3 \times 10^{-4 \sim -2} \lambda_{\rm eff} \kappa \sin\delta \,, \tag{3.12}$$

where κ is the washout factor and c_s is the conversion rate of the lepton number asymmetry into the baryon number asymmetry by the sphaleron interaction and takes a value of $c_s = O(1)$ [41]. To derive a last numerical factor, we used the value $|\varepsilon|/|\mu| = 10^{-4\sim-3}$ which has been referred in the footnote below eq. (3.5). If there is no entropy production after the decay of $\tilde{\Phi}_0$ condensate, this baryon number asymmetry is understood to represent the one of the present universe. Comparing this with the present predicted value $Y_B = (0.6 \; - \; 1)) \times 10^{-10}$, we find that $2 \times 10^{-9} \lesssim \lambda_{\rm eff} \kappa \sin\delta \lesssim 3 \times 10^{-7}$ should be satisfied. By substituting this into eq. (3.9), we obtain $2 \times 10^{-2} \lesssim \kappa \sin\delta \lesssim 1$. It is noticeable that this condition can be satisfied by the reasonable values of $\sin\delta$ and κ.

As we noted at the end of last section, however, there can be an additional reheating due to the decay of the $\tilde{\phi}$ condensate into the light SM fields. If this decay occurs after the decay of the $\tilde{\Phi}_0$ condensates, the entropy produced by this $\tilde{\phi}$ decay will dilute the baryon number asymmetry obtained from eq. (3.12) in such a way as [32]

$$Y_B \sim \frac{g_2^2 + g_1^2}{16\pi} c_s \kappa \frac{|\varepsilon|^2}{|\mu|^2} \sin\delta \frac{T_{\rm RH}^{\tilde{\phi}}}{M_{\rm susy}} \sim 3 \times 10^{-11 \sim -8} \kappa \sin\delta \frac{T_{\rm RH}^{\tilde{\phi}}}{M_{\rm susy}} \,. \tag{3.13}$$

[14] This condition may be in a delicate position in this model. For more precise analysis, we need to check it numerically. The ambiguity related to this point will be taken into account as the washout factor κ here.

It should be noted that this result is independent of the value of $\lambda_{\rm eff}$. The period of $\tilde{\phi}$ decay depends on the power r in W_2. In the case of $r = 5$ where $\phi QQQQ\bar{U}/\tilde{M}_{\rm pl}^3$ is, for example, contained in W_2,[15] the decay of $\tilde{\phi}$ occurs sufficiently after the one of $\tilde{\Phi}_0$ and eq. (2.20) shows that the reheating temperature is $T_{\rm RH}^{\tilde{\phi}} \sim O(10)$ GeV. In this case the required value of δ to realize the observed Y_B should be maximal and $\kappa \sim 1$. The $r \geq 6$ case seems to be ruled out based on the baryogenesis.

4. Summary

We have studied a D-term inflation scenario in a supersymmetric model extended from the MSSM by both the two abelian factor groups and several SM singlet chiral superfields. One of these groups is assumed to have a Fayet-Iliopoulos D-term ξ. Although the inflation scenario seems to be very similar to the ordinary D-term inflation model, our model can also have a lot of interesting features in particle physics. The condensates of the additionally introduced SM singlet fields to induce the inflation can be directly related to both the solution of the μ-problem and the explanation of the neutrino masses. They cause small bilinear R-parity violating terms which can induce the neutrino-neutralino mixing. The small neutrino masses can be generated through the weak scale seesaw mechanism based on this mixing. The decay of such condensates can also induce the leptogenesis through the effective coupling which violates the lepton number spontaneously. Although the reheating temperature due to these decays is not so high that the gravitino problem can be escapable, it is high enough that the electroweak sphaleron interaction can transfer the generated lepton number asymmetry into the baryon number asymmetry. This model can reconcile or relax the discrepancy between both values of ξ which are required by the CMB data and by the weakly-coupled superstring models if the SM singlet fields relevant to the inflation have appropriate charges of the additionally introduced abelian factor groups. Although there remains a lot of unstudied subjects, this kind of model seems to have interesting features in both astrophysics and particle physics.

Acknowledgments

This work is supported in part by a Grant-in-Aid for Scientific Research (C) from Japan Society for Promotion of Science (No. 11640267 and No. 14540251) and also by a Grant-in-Aid for Scientific Research on Priority Areas (A) from The Ministry of Education, Science, Sports and Culture (No. 14039205).

References

[1] G.F. Smoot et al., *Structure in the COBE DMR first year maps*, *Astrophys. J.* **396** (1992) L1; C.L. Bennett et al., *Four year COBE DMR cosmic microwave background observations: maps and basic results*, *Astrophys. J. Lett.* **464** (1996) L1.

[15] We can easily introduce a discrete symmetry to realize this by slightly modifying eq. (3.7).

[2] Super-Kamiokande collaboration, Y. Suzuki, *Solar neutrino results from Super-Kamiokande*, *Nucl. Phys.* **91** *(Proc. Suppl.)* (2001) 29.

[3] For a review see, D.H. Lyth and A. Riotto, *Particle physics models of inflation and the cosmological density perturbation*, *Phys. Rept.* **314** (1999) 1 [hep-ph/9807278].

[4] H. Murayama, H. Suzuki, T. Yanagida and J. Yokoyama, *Chaotic inflation and baryogenesis by right-handed sneutrinos*, *Phys. Rev. Lett.* **70** (1993) 1912.

[5] D. Suematsu and Y. Yamagishi, *Chaotic inflation based on an abelian D flat direction*, *Mod. Phys. Lett.* **A 10** (1995) 2923 [hep-ph/9504395].

[6] T. Gherghetta and G.L. Kane, *Chaotic inflation and a radiatively generated intermediate scale in the supersymmetric standard model*, *Phys. Lett.* **B 354** (1995) 300 [hep-ph/9504420].

[7] M. Bastero-Gil and S.F. King, *A next-to-minimal supersymmetric model of hybrid inflation*, *Phys. Lett.* **B 423** (1998) 27 [hep-ph/9709502];
G.L. Kane and S.F. King, *Intermediate scale supersymmetric inflation, matter and dark energy*, *New J. Phys.* **3** (2001) 21 [hep-ph/0109118].

[8] G. Lazarides, *Degenerate neutrinos and supersymmetric inflation*, *Phys. Lett.* **B 452** (1999) 227 [hep-ph/9812454].

[9] G.R. Dvali, G. Lazarides and Q. Shafi, *mu problem and hybrid inflation in supersymmetric* $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$, *Phys. Lett.* **B 424** (1998) 259 [hep-ph/9710314].

[10] K. Choi, E.J. Chun and J.E. Kim, *Cosmological implications of radiatively generated axion scale*, *Phys. Lett.* **B 403** (1997) 209 [hep-ph/9608222].

[11] G. Dvali and A. Riotto, *Minimalism in inflation model building*, *Phys. Lett.* **B 417** (1997) 20 [hep-ph/9706408].

[12] A.D. Linde, *Chaotic inflation*, *Phys. Lett.* **B 129** (1983) 177.

[13] A.D. Linde, *Eternal extended inflation and graceful exit from old inflation without jordan-brans-dicke*, *Phys. Lett.* **B 249** (1990) 18.

[14] D.H. Lyth and E.D. Stewart, *Cosmology with a TeV mass GUT Higgs*, *Phys. Rev. Lett.* **75** (1995) 201 [hep-ph/9502417]; *Thermal inflation and the moduli problem*, *Phys. Rev.* **D 53** (1996) 1784 [hep-ph/9510204].

[15] D.H. Lyth and E.D. Stewart, *More varieties of hybrid inflation*, *Phys. Rev.* **D 54** (1996) 7186 [hep-ph/9606412];
E. Halyo, *Hybrid inflation from supergravity D-terms*, *Phys. Lett.* **B 387** (1996) 43 [hep-ph/9606423].

[16] E.J. Copeland, A.R. Liddle, D.H. Lyth, E.D. Stewart and D. Wands, *False vacuum inflation with Einstein gravity*, *Phys. Rev.* **D 49** (1994) 6410 [astro-ph/9401011].

[17] E. Halyo, *Hybrid inflation from supergravity D-terms*, *Phys. Lett.* **B 387** (1996) 43 [hep-ph/9606423];
P. Binétruy and G.R. Dvali, *D-term inflation*, *Phys. Lett.* **B 388** (1996) 241 [hep-ph/9606342].

[18] G. Cleaver, M. Cvetič, J.R. Espinosa, L. Everett, P. Langacker and J. Wang, *Physics implications of flat directions in free fermionic superstring models, 1. Mass spectrum and couplings*, *Phys. Rev.* **D 59** (1999) 055005 [hep-ph/9807479]; *Physics implications of flat directions in free fermionic superstring models, 2. Renormalization group analysis*, *Phys. Rev.* **D 59** (1999) 115003 [hep-ph/9811355].

JHEP10(2002)014

[19] S. Coleman and E. Weinberg, *Radiative corrections as the origin of spontaneous symmetry breaking*, *Phys. Rev.* **D 7** (1973) 1888.

[20] J.R. Espinosa, A. Riotto and G.G. Ross, *D-term inflation in superstring theories*, *Nucl. Phys.* **B 531** (1998) 461 [hep-ph/9804214].

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

[22] M. Dine, N. Seiberg and E. Witten, *Fayet-Iliopoulos terms in string theory*, *Nucl. Phys.* **B 289** (1987) 589;
J.J. Atick, L.J. Dixon and A. Sen, *String calculation of Fayet-Iliopoulos d terms in arbitrary supersymmetric compactifications*, *Nucl. Phys.* **B 292** (1987) 109.

[23] A. Riotto, *D-term inflation: the good, the bad and the ugly*, hep-ph/9710329.

[24] D. Suematsu, *Vacuum structure of the μ-problem solvable extra* U(1) *models*, *Phys. Rev.* **D 59** (1999) 055017 [hep-ph/9808409].

[25] P. de Bernardis et al., *Multiple peaks in the angular power spectrum of the cosmic microwave background: significance and consequences for cosmology*, *Astrophys. J.* **564** (2002) 559 [astro-ph/0105296].

[26] G.R. Dvali, Q. Shafi and R. Schaefer, *Large scale structure and supersymmetric inflation without fine tuning*, *Phys. Rev. Lett.* **73** (1994) 1886 [hep-ph/9406319].

[27] J. García-Bellido and A.D. Linde, *Preheating in hybrid inflation*, *Phys. Rev.* **D 57** (1998) 6075 [hep-ph/9711360];
L. Kofman, A. Linde and A.A. Starobinsky, *Towards the theory of reheating after inflation*, *Phys. Rev.* **D 56** (1997) 3258 [hep-ph/9704452].

[28] D. Suematsu and Y. Yamagishi, *Radiative symmetry breaking in a supersymmetric model with an extra* U(1), *Int. J. Mod. Phys.* **A 10** (1995) 4521 [hep-ph/9411239].

[29] G. Cleaver, M. Cvetič, J.R. Espinosa, L.L. Everett and P. Langacker, *Intermediate scales, μ parameter and fermion masses from string models*, *Phys. Rev.* **D 57** (1998) 2701 [hep-ph/9705391].

[30] D. Suematsu and G. Zoupanos, *mu-term due to the non-universal supersymmetry breaking and the Higgs mass*, *J. High Energy Phys.* **06** (2001) 038 [hep-ph/0102096].

[31] J.R. Ellis, A.D. Linde and D.V. Nanopoulos, *Inflation can save the gravitino*, *Phys. Lett.* **B 118** (1982) 59.

[32] C.F. Kolda and J. March-Russell, *Supersymmetric D-term inflation, reheating and Affleck-Dine baryogenesis*, *Phys. Rev.* **D 60** (1999) 023504 [hep-ph/9802358].

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

[34] M. Gell-Mann, P. Ramond and S. Slansky, *Complex spinors and unified theories*, in *Supergravity*, P. van Nieuwenhuizen and D. Freedman eds., North-Holland, 1979, p. 315;
T. Yanagida, in *Proceedings of the workshop on unified theory and baryon number in the universe*, A. Sawada and H. Sugawara eds., KEK, 1979.

[35] For example see, R. Barbier et al., *Report of the group on the R-parity violation*, hep-ph/9810232.

JHEP10(2002)014

[36] F. de Campos, M.A. García-Jareño, A.S. Joshipura, J. Rosiek and J.W.F. Valle, *Novel scalar boson decays in SUSY with broken R-parity*, *Nucl. Phys.* **B 451** (1995) 3 [hep-ph/9502237]; T. Banks, Y. Grossman, E. Nardi and Y. Nir, *Supersymmetry without R-parity and without lepton number*, *Phys. Rev.* **D 52** (1995) 5319 [hep-ph/9505248]; A.S. Joshipura and M. Nowakowski, *'Just so' oscillations in supersymmetric standard model*, *Phys. Rev.* **D 51** (1995) 2421 [hep-ph/9408224].

[37] M.A. Díaz, J.C. Romão and J.W.F. Valle, *Minimal supergravity with R-parity breaking*, *Nucl. Phys.* **B 524** (1998) 23 [hep-ph/9706315]; J.W.F. Valle, *Supergravity unification with bilinear R-parity violation*, in *Proceedings 6th international symposium on particles, strings and cosmology*, (PASCOS '98), Boston 1998, p. 502 and references therein.

[38] D. Suematsu, *Neutrino mass due to the mixing among neutrinos and gauginos*, *Phys. Lett.* **B 506** (2001) 131 [hep-ph/0012321]; *The origin of quark and lepton mixings*, *Phys. Rev.* **D 64** (2001) 073013 [hep-ph/0104187]; J. Kubo and D. Suematsu, *Suppressing the μ and neutrino masses by a superconformal force*, *Phys. Rev.* **D 64** (2001) 115014 [hep-ph/0107133].

[39] M. Hirsch, M.A. Díaz, W. Porod, J.C. Romão and J.W.F. Valle, *Neutrino masses and mixings from supersymmetry with bilinear R-parity violation: a theory for solar and atmospheric neutrino oscillations*, *Phys. Rev.* **D 62** (2000) 113008 [hep-ph/0004115].

[40] D. Suematsu, *μ-term as the origin of lepton number asymmetry*, hep-ph/0207039.

[41] S.Y. Khlebnikov and M.E. Shaposhnikov, *The statistical theory of anomalous fermion number nonconservation*, *Nucl. Phys.* **B 308** (1988) 885; J.A. Harvey and M.S. Turner, *Cosmological baryon and lepton number in the presence of electroweak fermion number violation*, *Phys. Rev.* **D 42** (1990) 3344.

JHEP10(2002)014