Dark world and baryon asymmetry from a common source

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Abstract. We study generation of baryon number asymmetry and both abundance of dark matter and dark energy on the basis of global symmetry and its associating flat directions in a supersymmetric model. We assume the existence of a model independent axion, which is generally expected in the effective superstring theory. If we consider a combined field of the model independent axion and a pseudo-Nambu–Goldstone boson coming from spontaneous breaking of the global symmetry, its potential can be sufficiently flat and then it may present a candidate for the dark energy as a quintessential axion. Both the baryon asymmetry and the dark matter are supposed to be produced nonthermally as the asymmetry of another global charge through the Affleck–Dine mechanism along the relevant flat direction. Its decay to the observable and hidden sectors explains the baryon number asymmetry and the dark matter abundance, respectively.

Keywords: dark matter, dark energy theory, baryon asymmetry, axions
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1. Introduction

Recent observations of the anisotropy of cosmic microwave background [1], the distance–redshift relation of type Ia supernovae [2] and the large scale structure [3] show that there is a large amount of energy, called dark energy, which has negative pressure to cause acceleration of the expansion of the universe. Analyses of these data also allow us to determine energy density parameters with good accuracy. This situation for these astrophysical observations makes us confront several difficult problems, that is, what is the origin of the dark energy, why each energy density parameter takes the present values and so on. In particular, since each energy density is expected to have a different origin and also behaves in a very different manner during the expansion of the universe, it seems very mysterious why they take similar order values in the present universe. The investigation of these problems might give us some clues to access new physics beyond the standard model of particle physics. In the supersymmetric framework, in particular, many works have been done, mainly related to the thermal neutralino dark matter [4].

As one of the interesting approaches to these problems, one may consider them on the basis of some global symmetry. As proposed in [5], both abundance of the baryons and the dark matter may be understood from a viewpoint of the charge asymmetry of a suitable common global symmetry. Although this idea seems to be very elegant and also able to explain naturally why they have the similar order values in the present universe, realistic models in this direction seem not to have been constructed much in the context of a supersymmetric framework.

The Lagrangian of the minimal supersymmetric standard model (MSSM) has only baryon number ($B$) and lepton number ($L$) as its global symmetry. Even if we put aside

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1 An example of this kind of recent works may be found in [6]. We can also find several other attempts to explain both the dark matter and baryon asymmetry based on a single source. See, for example, [7,8]. We should note that the tuning of the dark matter mass to the proton mass can remain as an unsolved problem even in the case where both number density of the baryons and the dark matter are related to each other through a certain symmetry.
the dark energy, it seems to be difficult to explain both abundance of the baryons and the dark matter on the basis of these global symmetries. However, if we introduce new singlet chiral superfields in the MSSM, the global symmetry can be extended and then the idea given in [5] may be applicable to the supersymmetric model by using those symmetries. In the previous paper [9], we have discussed such a possibility to explain both abundance of the baryons and the dark matter on the basis of a common global symmetry. The global charge asymmetry induced by the Affleck–Dine mechanism [10] in a $D$-flat direction associated with the introduced singlet fields is liberated into the observable and hidden sectors. These decay products are shown to explain the baryon asymmetry and the dark matter abundance, respectively. Although we left the dark energy untouched in that study, it seems interesting to investigate a new possibility, that the dark energy may be explained in the same framework.

For the dark energy, many works have been done by assuming the existence of a slowly rolling homogeneous scalar field with extremely small mass and a large vacuum expectation value, which is called quintessence [11]–[13]. Although they may be able to explain the observational data to some level, many of them seem not to be motivated from particle physics. Moreover, if we try to consider the quintessence in the supersymmetric framework, it is feared that the supergravity corrections violate the slow roll nature of the quintessence as in the case of the inflaton [14]. To escape this problem, we may impose some symmetry which can protect the flatness of the potential from such corrections. As such a solution, one may consider that the quintessence field is a pseudo-Nambu–Goldstone (pNG) boson of some global symmetry [15]. In this case, the flatness of its potential is guaranteed by the shift symmetry of the pNG field. Thus, the axion type field seems to be a promising candidate for the dark energy, which can also be supported by a suitable motivation from particle physics. In such a scenario, a crucial problem is how we can realize the tiny mass scale suitable for the dark energy. On this point, there is an interesting proposal in which a model independent axion in the superstring is used as a candidate for the quintessence [16].

In this paper we try to extend the idea given in [16] so that the scenario for the baryon asymmetry and the dark matter abundance given in [9] is embedded into the same framework. If it is possible, we may have a common source for the baryon asymmetry, the dark matter and the dark energy in the flat directions associated with the extended global symmetry motivated by particle physics. The remaining parts of this paper are organized as follows. In section 2 we present a supersymmetric model which is mainly characterized by the structure of the hidden sector. Several features of the model are discussed. In section 3 the dark energy is studied in this model. We show that two axion fields associated with the almost flat directions can have hierarchical mass and they can play the role of the quintessence and the invisible axion for the strong CP problem, respectively. In section 4 we discuss how both abundance of the baryon and the dark matter can be produced in this framework. Section 5 is devoted to the summary. Some details of the calculation are explained in two appendices.

2. A model with $D$-flat directions

We consider a supersymmetric model with observable and hidden sectors. The observable sector is supposed to be composed of the MSSM contents and the MSSM singlet chiral
superfields $\hat{N}_j$, $\hat{S}_1$. The superpotential is written as

$$W_{ob} = \sum_{i,j=1}^{3} \left( h^{ij}_u \hat{Q}_i \hat{U}_j \hat{H}_2 + h^{ij}_d \hat{Q}_i \hat{D}_j \hat{H}_1 + h^{ij}_e \hat{L}_i \hat{E}_j \hat{H}_1 + h^{ij}_s \hat{L}_i \hat{N}_j \hat{H}_2 \right) + \lambda_1 \hat{S}_1 \hat{H}_1 \hat{H}_2, \quad (1)$$

which is similar to the usual next MSSM superpotential up to a cubic term of $\hat{S}_1$. The hidden sector is assumed to be composed of two parts. One is a super-Yang–Mills part composed of an $SU(N)$ vector superfield and three sets of chiral superfields such as $\hat{F}^\alpha \left( N, \alpha = 1, \ldots , n_{G_1} + n_{G_2} \right)$, $\hat{G}^\beta_1 \left( N, \beta = 1, \ldots , n_{G_1} \right)$ and $\hat{G}^\gamma_2 \left( \bar{N}, \gamma = 1, \ldots , n_{G_2} \right)$, where $N$ and $\bar{N}$ stand for the fundamental representation of $SU(N)$ and its conjugate representation, respectively. We assume that the supersymmetry is broken in this part of the hidden sector. Another part is composed of several gauge singlet chiral superfields $\hat{S}_2$, $\hat{C}_i (i = 1, \ldots , 5)$, $\hat{\phi}$ and $\hat{\psi}$.

The superpotential for these chiral superfields is assumed to be

$$W_{hid}^{YM} = \sum_{\beta=1}^{n_{G_1}} \frac{g_1}{M_{pl}} \hat{\phi} \hat{F}^\beta \hat{G}^\beta_1 + \sum_{\gamma=1}^{n_{G_2}} \frac{g_2}{M_{pl}} \hat{\psi} \hat{G}^\gamma \hat{G}^\gamma_2,$$

$$W_{hid}^S = \lambda_2 \hat{S}_2 \hat{C}_1 \hat{C}_2 + \frac{f_1}{M_{pl}} \hat{\phi}^2 \hat{C}_1 \hat{C}_3 + \frac{f_2}{M_{pl}} \hat{\phi}^2 \hat{C}_2 \hat{C}_4 + \frac{f_3}{M_{pl}} \hat{\phi}^2 \hat{S}_2 \hat{C}_5, \quad (2)$$

where $M_{pl}$ is the reduced Planck mass and all coefficients of these operators are supposed to be real and $O(1)$. The observable and hidden sectors are assumed to be connected only through the gravitational interactions which are represented by the following superpotential:

$$W_{mix} = \frac{d_1}{M_{pl}^2} \hat{S}_1 \hat{S}_2 + \sum_{\gamma=1}^{n_{G_2}} \frac{d_2}{M_{pl}} \hat{F}^\gamma \hat{G}^\gamma_2 \hat{H}_1 \hat{H}_2, \quad (3)$$

where the coefficients $d_{1,2}$ are also assumed to be real and $O(1)$. The operators in $W_{mix}$ play essential roles in the present scenario as discussed below.

Many singlet chiral superfields are introduced in the hidden sector. This seems to be unavoidable to realize several physical results simultaneously, keeping the global symmetry in the hidden sector which is essential for the present approach. Although the detailed discussions will be presented in the following, it may be useful to present a brief summary of their specific roles here. The global charge asymmetry related to both the dark matter and baryon asymmetry is generated through the flat direction of scalar components of $\hat{S}_{1,2}$. $\hat{\phi}$ contributes to the supersymmetry breaking in the observable sector through its coupling with $\hat{F}^\alpha$ and $\hat{G}^\gamma_2$ in $W_{hid}^{YM}$. The vacuum expectation value of the scalar component of $\hat{\phi}$ generates the mass of dark matter candidates $\hat{C}_i$ without violating the global symmetry through the couplings in $W_{hid}^S$. The invariance of these couplings is to be unavoidable to realize several physical results simultaneously, keeping the global symmetry in the hidden sector which is essential for the present approach. Although the detailed discussions will be presented in the following, it may be useful to present a brief summary of their specific roles here. The global charge asymmetry related to both the dark matter and baryon asymmetry is generated through the flat direction of scalar components of $\hat{S}_{1,2}$. $\hat{\phi}$ contributes to the supersymmetry breaking in the observable sector through its coupling with $\hat{F}^\alpha$ and $\hat{G}^\gamma_2$ in $W_{hid}^{YM}$. The vacuum expectation value of the scalar component of $\hat{\phi}$ generates the mass of dark matter candidates $\hat{C}_i$ without violating the global symmetry through the couplings in $W_{hid}^S$. The invariance of these couplings

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\[ \text{In this paper, we put a hat on the character for the chiral superfields in both sectors. For the scalar component of the hidden sector chiral superfield, we put a tilde on it. For its fermionic component, we use the same character as the superfield without the hat.} \]

\[ \text{This part is relevant to the explanation of the dark matter abundance as seen below. Although it may be possible to consider similar structure to the observable sector for this part, we assume this simple structure here, for simplicity.} \]
Table 1. A charge assignment for the chiral superfields in the observable sector and the anomaly $A_3$ and $A_2$ for $SU(3)$ and $SU(2)$. A trivial $Z_3$ charge is assigned for all fields.

| $Z_{18}$ | $Q_{FQ}$ | $Q_R$ | $B$ | $X$ | $L$ | $A_3$ | $A_2$ |
|---------|---------|-------|-----|-----|-----|-------|-------|
| 1       | 0       | 1     | 0   | 0   | 0   | 0     | 1     |
| 5       | 0       | 1     | 1   | 0   | 0   | 1     | 0     |
| 5       | 0       | 1     | 0   | 0   | 0   | 0     | 0     |
| 1       | 0       | 0     | 0   | 0   | 0   | 0     | 0     |
| 5       | 0       | 0     | 0   | 0   | 0   | 0     | 0     |
| 2       | 0       | 1     | 0   | 0   | 0   | 0     | 0     |
| 3       | 0       | 1     | 0   | 0   | 0   | 0     | 0     |
| 1       | 0       | 0     | 0   | 0   | 0   | 0     | 0     |
| 0       | 0       | 0     | 0   | 0   | 0   | 0     | 0     |

Table 2. A charge assignment for the chiral superfields in the hidden sector and the anomaly $A_N$ for $SU(N)$. All fields have $B = L = 0$.

| $F$ | $G_1$ | $G_2$ | $S_2$ | $C_1$ | $C_2$ | $C_3$ | $C_4$ | $C_5$ | $\phi$ | $\psi$ | $A_N$ |
|-----|------|------|------|------|------|------|------|------|-------|-------|-------|
| 2   | 0    | 1    | 0    | 1    | 2    | 1    | 0    | 2    | 0     | 0     | $2/3(2n_{G_1} + 3n_{G_2})$ |
| 3   | 0    | 0    | 0    | 0    | 0    | 0    | 0    | 0    | 0     | 0     | $n_{G_2} - N$ |
| 1   | 0    | 0    | 0    | 0    | 0    | 0    | 0    | 0    | 0     | 0     | $n_{G_2} - N$ |
| 0   | 0    | 0    | 0    | 0    | 0    | 0    | 0    | 0    | 0     | 0     | $3/2n_{G_2} - N$ |

under the global symmetry requires us to introduce a rather large number of $\hat{C}_i$. $\hat{\phi}$ also contributes to the generation of the dark energy through the coupling in $W_{\text{YM}}$.

In order to control the structure of $W_{\text{hid}}$ and $W_{\text{mix}}$, we may impose discrete symmetry $Z_{18} \times Z_3$. We assign their charges for each chiral superfield in the observable and hidden sectors as shown in tables 1 and 2. Each operator in $W_{\text{ob}}$, $W_{\text{hid}}$ and $W_{\text{mix}}$ is $Z_{18} \times Z_3$ invariant. The operators in $W_{\text{hid}}$ and $W_{\text{mix}}$ are the lowest order ones which can be constructed in a $Z_{18} \times Z_3$ invariant way by using the fields in both sectors. It can be checked that this $Z_{18} \times Z_3$ has no anomaly with respect to both the SM gauge group and $SU(N)$ if we take $n_{G_1} = 3$ and $n_{G_2} = 1$.

In the MSSM, both operators $\hat{H}_1\hat{H}_2$ and $\hat{L}_i\hat{H}_2$ are gauge invariant. Thus, as shown in $W_{\text{ob}}$, we can construct two gauge invariant dimension three operators $\hat{S}_i\hat{H}_1\hat{H}_2$ and $\hat{L}_i\hat{N}_j\hat{H}_2$ by introducing the singlet chiral superfields $\hat{S}_i$ and $\hat{N}_j$. If we add $\hat{S}_i\hat{H}_1\hat{H}_2$ to the MSSM superpotential and remove the ordinary $\mu$-term as in $W_{\text{ob}}$, the global symmetry is generally extended. In fact, in the present model two new global Abelian symmetries

\footnote{Among gauge invariant renormalizable operators constructed from the contents given in table 1 and 2, we cannot exclude $\hat{S}_i^3$ by imposing $Z_{18} \times Z_3$ only. We need an additional symmetry like $Z_2$ to forbid it. Although its charge can be consistently assigned to all fields, we do not present it explicitly here and only assume implicitly that such a symmetry is imposed.}
appear in addition to the $B$ and $L$ symmetries as long as we restrict our target to the renormalizable operators \cite{17,18,9}. They are taken as the Peccei–Quinn type symmetry and the $R$ symmetry\footnote{Since the ordinary Peccei–Quinn symmetry has no $SU(2)$ gauge anomaly, the present one is not such a symmetry in an exact sense.}. We present an example of their charge assignment for each chiral superfield in the observable and hidden sectors in tables 1 and 2, respectively. Gauginos are considered to have $Q_R = -1$ only. We note that the global symmetries in the hidden and observable sectors can be related to each other since the model is assumed to have the operators in $W_{\text{mix}}$. This charge assignment will be used in the following discussion.

Among these four global Abelian symmetries, $U(1)_{B-L}$ and $U(1)_X$ remain as those with no $SU(3)$ and $SU(2)$ gauge anomalies in the observable sector. The $U(1)_X$ symmetry is a kind of $Y$-symmetry and its charge can be taken as the linear combination of the four global Abelian charges in such a way as $X = \frac{1}{3}(B+L-10Q_{\text{PQ}})+Q_R$. In the hidden sector, there is also an $SU(N)$ anomaly free global Abelian symmetry in the case of $n_{G_1} = 3$, $n_{G_2} = 1$ and $N = 5$. This charge can be represented as $Y = X - 8Q_{\text{PQ}}$. These charges are listed in tables 1 and 2. It may be interesting to note that the same values of $n_{G_1}$ and $n_{G_2}$ make $Z_{18} \times Z_3$ gauge anomaly free as mentioned before.

It is useful to note that only the soft supersymmetry breaking operator associated with the first term in $W_{\text{mix}}$ can violate the $X$ conservation in the Lagrangian derived from $W_{\text{ob}}, W_{\text{hid}}$ and $W_{\text{mix}}$. As we will see later, it can play an essential role to generate the $X$ asymmetry through the AD mechanism in the $D$-flat direction of the scalar components of $\hat{\phi}_{1,2}$.

Using $U(1)_{\text{PQ}}$ and $U(1)_{R}$, we can also construct a global Abelian symmetry $U(1)_{X}$, which is $SU(2)$ anomaly free, but have $SU(3)$ and $SU(N)$ anomalies. As such an example, we can consider $X = -\frac{42}{3}Q_{\text{PQ}} + Q_R$. In the next section, we will suggest that this global symmetry can be related to the appearance of a quintessential axion.

Here we make some assumptions for the dynamics in the super-Yang–Mills part of the hidden sector along the line of the arguments given in \cite{16}. The $SU(N)$ gaugino $\lambda_h$ and the scalar components $\tilde{F}, \tilde{G}_{1,2}$ of the chiral superfields $\tilde{F}, G_{1,2}$ are supposed to condense at the scale $\Lambda_h$. The condensations of the gaugino, the scalar components $\tilde{F}$ and $\tilde{G}_{1,2}$ are defined by

\begin{equation}
\langle \lambda_h \lambda_h \rangle \equiv \Lambda_h^3 e^{i(\theta_{\gamma,0}/\Lambda_h)}, \quad \langle \tilde{F}^\beta \tilde{G}_{1,2}^\beta \rangle \equiv \Lambda_h^2 e^{i(\theta_{G_{1,2}}/\Lambda_h)}, \quad \langle \tilde{F}^\gamma \tilde{G}_{2}^\gamma \rangle \equiv \Lambda_h^2 e^{i(\theta_{G_{1,2}}/\Lambda_h)},
\end{equation}

where $\beta = 1, \ldots, n_{G_1}$ and $\gamma = 1, \ldots, n_{G_2}$. Although these condensations do not break the supersymmetry by themselves, the $U(1)_X$ and $U(1)_{X}$ symmetries are spontaneously broken in the hidden sector by these condensations\footnote{It should be noted that these breakings have no substantial effect on the singlet part of the hidden sector, which will be discussed later.} and the NG bosons $\theta_{G_{1,2}}(x)$ and $\theta_{\hat{\phi}}(x)$ appear as composite states. Since $U(1)_X$ has $SU(3)$ and $SU(N)$ anomalies, the corresponding NG boson gets a mass through these effects to be the pNG boson\footnote{The scenario for the dark energy in \cite{16} is regarded to be constructed from this pNG and the model independent axion $a_{\text{MT}}(x)$.}. It can be expressed by using $\theta_{G_1}, \theta_{G_2}$ and $\theta_{\lambda}$ as [19, 20]

\begin{equation}
\theta_{\hat{\phi}} = \frac{2n_{G_2}}{(4n_{G_2}^2 + 25n_{G_2}^2 + 4)^{1/2}}.
\end{equation}
The supersymmetry breaking in the observable sector is expected to appear indirectly through the gravitational interaction which mediates the effects of the condensations (4) \cite{21}. In the general supergravity framework, the order parameter for the supersymmetry breaking in the observable sector is expressed as

$$F = \sum_i F^i + F_A + \cdots = \sum_i \frac{\partial W_{\text{hid}}}{M_{\text{pl}}} \xi^{K/2M_{\text{pl}}^2} + \frac{1}{4M_{\text{pl}}} \langle \lambda_h \lambda_h \rangle + \cdots, \quad (6)$$

where $K$ is the Kähler potential and the minimal kinetic terms are assumed. Since the condensations (4) are assumed to be generated in the hidden sector, the dominant supersymmetry breaking effect in the observable sector is expected to come from the first term of equation (6). In fact, the second term of $W_{\text{hid}}^{\text{YM}}$ generates the dominant effect of the condensation mediated into the observable sector\footnote{The corresponding term to this does not exist in the model of \cite{16} and the difference from our scenario is caused by this point. Although there may be other contributions to the supersymmetry breaking such as the one due to the modulus $F$ term, for example, we do not refer to it here.}. This supersymmetry breaking effect is expected to have a typical scale

$$M_{\text{susy}} \equiv \sqrt{\langle F^\dagger F \rangle} = O \left( \frac{\Lambda_h^2}{M_{\text{pl}}} \right). \quad (7)$$

If we require $M_{\text{susy}} = O(10^3)$ GeV, $\Lambda_h$ should be an intermediate scale such as $\Lambda_h = O(10^{10})$ GeV.

Since our model has the singlet chiral superfields $\tilde{S}_{1,2}$ which have $U(1)_{\tilde{X}}$ charges, they may give additional contributions to the above mentioned pNG boson if their scalar components get the vacuum expectation values to induce the spontaneous symmetry breaking of $U(1)_{\tilde{X}}$. We parametrize the fluctuation of the scalar components $\tilde{S}_{1,2}$ from their vacuum values as

$$\tilde{S}_1 \equiv (u_1 + \eta_1(x)) e^{i \theta_1(x)/u_1}, \quad \tilde{S}_2 \equiv (u_2 + \eta_2(x)) e^{i \theta_2(x)/u_2}, \quad (8)$$

where $u_i = |\langle \tilde{S}_i \rangle|$. Originally, there are two degrees of freedom for the phases of $\tilde{S}_{1,2}$. However, one of their linear combination $\xi \equiv (2\sqrt{3}\theta_1 + 3\sqrt{2}\theta_2)/5$ can be fixed dynamically to a constant value as a potential minimum during the early stage of the evolution of $\tilde{S}_{1,2}$. We will discuss this point later. Thus, only its orthogonal state $\theta \equiv (-3\sqrt{2}\theta_1 + 2\sqrt{3}\theta_2)/5$ can contribute to the discussing pNG boson $\theta_{\tilde{X}}$.

We make an additional assumption related to this pNG sector, which is natural if the model is supposed to be a low energy effective theory of some superstring. That is, the model is assumed to have a model independent axion $a_{\text{MI}}$. Its physical property can be briefly summarized as follows. It is a dual of the field strength $H_{\mu\nu\rho}$ of the second rank antisymmetric tensor $B_{\mu\nu}$ and is defined by $\partial^\mu a_{\text{MI}} = \epsilon^{\mu\nu\rho\sigma} H_{\nu\rho\sigma} \ [22]$. Its decay constant $f_{a_{\text{MI}}}$ is generally considered to be $f_{a_{\text{MI}}} = O(M_{\text{pl}})$. It constitutes an imaginary part of the dilaton and has universal effective couplings with non-Abelian gauge fields such as $(1/32\pi^2)(a_{\text{MI}}/f_{a_{\text{MI}}}) F^\alpha_{\mu\nu} \tilde{F}^{\alpha\mu\nu}$, where $F^\alpha_{\mu\nu}$ is the non-Abelian gauge field strength and $\tilde{F}^{\alpha\mu\nu}$ is its dual. In the present model, $a_{\text{MI}}$ has this type of coupling with the $SU(3)$ and $SU(N)$ gauge fields.
3. Dark energy candidate

In this section we show that our model can include a quintessential axion as a dark energy candidate in a similar way to [16]. The model is supposed to have two axions as discussed in the previous section. One is the pNG boson associated with the spontaneous breaking of $U(1)_X$ and the other one is the model independent axion. Since the operators in $W_{\text{obs}}$, $W_{\text{hid}}$ and $W_{\text{mix}}$ except for the first term in $W_{\text{mix}}$ are $U(1)_X$ invariant, the potential for these pseudo-scalar fields is expected to be induced mainly from their effective interactions with the $SU(3)$ and $SU(N)$ instantons. These interactions can be written as

$$\frac{2}{32\pi^2} \left( \frac{a}{f_a} + \frac{a_{\text{MI}}}{f_{\text{aMI}}} \right) F_{\mu \nu} \tilde{F}^{\mu \nu}, \quad \frac{2}{32\pi^2} \left( \frac{a_h}{f_h} + \frac{a_{\text{MI}}}{f_{\text{aMI}}} \right) G_{\mu \nu} \tilde{G}^{\mu \nu},$$

where $F_{\mu \nu}$ and $G_{\mu \nu}$ are the field strengths of $SU(3)$ and $SU(N)$ and their gauge indices are abbreviated. Axion fields $a$ and $a_h$ and their decay constants $f_a$ and $f_h$ are defined as [19]

$$a = \frac{5}{f_a} (u_1 \theta_1 + n_{G_2} \Lambda_h \theta_{G_2}) \simeq \frac{5n_{G_2} \Lambda_h}{f_a} \theta_{\tilde{X}} + \cdots,$$

$$f_a = 5 \left( u_1^2 + n_{G_2}^2 \Lambda_h^2 \right)^{1/2} \simeq 5n_{G_2} \Lambda_h,$$

$$a_h = \frac{1}{f_h} (5n_{G_2} \Lambda_h \theta_{G_2} + 2n_{G_1} \Lambda_h \theta_{G_1} - 2\Lambda_h \theta_{\lambda}) \simeq \frac{5n_{G_2} \Lambda_h}{f_h \gamma} \theta_{\tilde{X}},$$

$$f_h = \left( 25n_{G_2}^2 \Lambda_h^2 + 4n_{G_1}^2 \Lambda_h^2 + 4\Lambda_h^2 \right)^{1/2} \simeq \frac{5n_{G_2} \Lambda_h}{\gamma},$$

where $\gamma = 1 + ((4n_{G_1}^2 + 4/25n_{G_2}^2))^{-1/2}$. The ellipses in the expression of $a$ stand for the components orthogonal to $\theta_{\tilde{X}}$. In these formulae, we assume $|\langle H_{1,2} \rangle| \ll u_{1,2} \ll \Lambda_h$ in the case of $u_{1,2} \neq 0$. Then the contributions from the $U(1)_X$ breaking due to $\langle H_{1,2} \rangle$ and $u_{1,2}$ are neglected.

If we assume that the $SU(N)$ gaugino and the fermionic components of the chiral superfields $\tilde{F}$ and $\tilde{G}_{1,2}$ have the masses $M_{\tilde{g}}$, $m_{G_1}$ and $m_{G_2}$, the scalar potential for these axions generated through the interactions (9) may be written as [16]

$$V = m_{\text{pl}}^2 f_{\pi}^2 \left[ 1 - \cos \left( \frac{f_h}{f_{\text{aMI}}} \gamma \right) \right] + \left[ 1 - \cos \left( \frac{f_h}{f_{\text{aMI}}} \gamma \right) \right] \left| \langle H_{1,2} \rangle \right| \simeq m_{\text{pl}}^2 f_{\pi}^2 \left[ 1 - \cos \left( \frac{f_h}{f_{\text{aMI}}} \gamma \right) \right] \left| \langle H_{1,2} \rangle \right| \simeq 10^{10} \text{ GeV},$$

where $V_h = m_{G_1}^4 m_{G_2}^4 M_{\text{pl}}^2 \Lambda_h^{4-n_{G_1}-n_{G_2}=N}$ and we take $f_{\text{aMI}} = M_{\text{pl}}$. The cosmological constant is assumed to be zero. For the observable sector, we use the ordinary invisible axion potential due to the QCD instanton [23]. One may wonder whether the higher order $\tilde{X}$ violating operators ignored in $W_{\text{hid}}$ and $W_{\text{mix}}$ might give a larger contribution to the $a_h$ potential compared with that given in equation (11). The most dangerous operator among the $\tilde{X}$ breaking higher dimensional ones is the supersymmetry breaking operator of the type $B/M_{\text{pl}}^{-3}(\tilde{F}_1 \cdots \tilde{F}_k)$ which corresponds to the superpotential operator $1/M_{\text{pl}}^{-3}(\tilde{F}_1 \cdots \tilde{F}_k)$. If we take $B \simeq 10^3$ GeV and $\langle \tilde{F} \rangle \simeq 10^{10}$ GeV as their conservative values, these effects cannot dominate the second term of equation (11). In fact, they determine the mass of $a_h$ only if $k \leq 5$ is satisfied. However, we can easily check that this
is not the case in the present model. Thus, it is sufficient to consider equation (11) as the dominant potential for \( a_h \).

Using the scalar potential (11), the mass matrix for two axions can be written in the basis \( (\theta_\chi, a_{\text{MI}}) \) as

\[
\begin{pmatrix}
\frac{1}{f_h}(\gamma^2 m_\pi^2 f_\pi^2 + V_h) & \frac{1}{f_h M_{\text{pl}}} (\gamma^2 m_\pi^2 f_\pi^2 + V_h) \\
\frac{1}{f_h M_{\text{pl}}} (\gamma^2 m_\pi^2 f_\pi^2 + V_h) & \frac{1}{f_h M_{\text{pl}}} (\gamma^2 m_\pi^2 f_\pi^2 + V_h)
\end{pmatrix}.
\]

Since \( \gamma \neq 1 \) is satisfied in the present model, we find that both eigenvalues are nonzero and they are expressed as

\[
m_{\text{QCD}}^2 \simeq \frac{m_\pi^2 f_\pi^2}{(f_h/\gamma)^2}, \quad m_{\text{aq}}^2 \simeq \frac{m_{G1} m_{G2} M_{\text{pl}} N^4}{M_{\text{pl}}^2} \frac{1}{M^2}.
\]

The corresponding eigenstates can be given by

\[
a_{\text{QCD}} \simeq \theta_\chi - \frac{f_h/\gamma}{M_{\text{pl}}} a_{\text{MI}}, \quad a_q \simeq \frac{f_h/\gamma}{M_{\text{pl}}} \theta_\chi + a_{\text{MI}}.
\]

Since the QCD axion \( a_{\text{QCD}} \) is dominated by \( \theta_\chi \), and its decay constant is given by \( f_h/\gamma \), it should satisfy the astrophysical constraint \( f_h/\gamma \simeq 10^{9-12} \) GeV [24]. This is found to be realized for the supposed values of \( \Lambda_h \) and \( \gamma \) here. Thus, \( a_{\text{QCD}} \) is expected to behave as the invisible axion to solve the strong CP problem.

It should be noted that both the origin of the \( \mu \) term and the candidate for the cold dark matter are intimately related to the scale of \( \Lambda_h \). In the present model, we have two sources for the \( \mu \) term, which are \( \mu = \lambda_1 (\langle \tilde{S}_1 \rangle) \) and \( \mu = (d_2/M_{\text{pl}}) (\langle \tilde{F} \tilde{G}_2 \rangle) \). For the first case, we should take account of the assumed relation \( \langle \langle \tilde{S}_1 \rangle \rangle \ll \Lambda_h \). As long as this is satisfied, \( \mu \) can take various values depending on the values of \( \langle \tilde{S}_1 \rangle \) and \( \lambda_1 \). On the other hand, for the latter case, if we take \( \Lambda_h \simeq 10^{12} \) GeV so that \( a_{\text{QCD}} \) can be a cold dark matter candidate [16], the \( \mu \) scale takes a large value such as \( O(10^6) \) GeV. In order to make its contribution to the \( \mu \) term suitable, we should assume \( \Lambda_h \sim 10^{10} \) GeV, which is just a supposed value here. Since \( f_h/\gamma \) seems to be too small for \( a_{\text{QCD}} \) to be the dark matter candidate in this case [25], we need another candidate for it. We will discuss this problem in the next section.

The formula in equation (13) shows that the mass of \( a_q \) can be strongly suppressed so that the present Hubble constant \( H_0 \) can be larger than this. Then, \( a_q \) can still now stay near the initial value at the beginning of the universe. It is naturally expected that \( a_q \) is far from the true vacuum. This means that the state equation for \( a_q \) can satisfy

\[
w = \frac{(1/2) a_q^2 - V(a_q)}{(1/2) a_q^2 + V(a_q)} \simeq -1
\]

as long as \( a_q \) takes a large value. Since \( a_q \) is dominated by \( a_{\text{MI}} \), it is natural to take the initial value of \( a_q \) as \( a_q \simeq M_{\text{pl}} \). If the energy density of \( a_q \) dominates that of the present universe, \( a_q \) controls the expansion behaviour of the universe. In this case the observed accelerated expansion of the present universe may be explained. If we assume \( a_q \simeq M_{\text{pl}} \),

\[\text{We note that the } \tilde{X} \text{ violating but } Z_{18} \times Z_3 \text{ invariant higher dimensional operators which include gauginos cannot give a dominant contribution either.}\]
the energy density stored in $a_q$ can be estimated as

$$\rho_{a_q} \simeq m_{G_1}^{-n_{G_1}} m_{G_2}^{-n_{G_2}} M_\Lambda^{-N} \Lambda_h^{4-n_{G_1} - n_{G_2} - N},$$

(16)

where we use $m_{a_q}^2$ in equation (13).

Here we assume that the $SU(N)$ gauginos get the supersymmetry breaking mass $M_\tilde{g} \simeq 10^3$ GeV. The fermionic components of $\hat{F}$ and $\hat{G}_{1,2}$ can get the masses through the couplings in $W_{\text{hid}}$ and $W_{\text{mix}}$ if we assume $\langle \tilde{\phi} \rangle \neq 0$ and $\langle \tilde{\psi} \rangle = 0$. They are expressed as

$$m_{G_1} = g_1 \frac{|\langle \tilde{\phi} \rangle|^2}{M_{pl}}, \quad m_{G_2} = d_2 \frac{|\langle H_1 \rangle \langle H_2 \rangle|}{M_{pl}}.$$  

(17)

The second formula results in $m_{G_2} \simeq 10^{-15}$ GeV. On the other hand, if we suppose $\langle \tilde{\phi} \rangle = O(10^{10})$ GeV, $m_{G_1} \sim 10$ GeV is induced. If we use these values in equation (16), the energy density of $a_q$ can be estimated as

$$\rho_{a_q} \simeq 10^{40 - 9n_{G_1} - 24n_{G_2} - 7N} \text{ GeV}^4.$$  

(18)

If we use $n_{G_1} = 3$, $n_{G_2} = 1$ and $N = 5$, which make the present model consistent, as seen in the previous section, we have $\rho_{a_q} \simeq 10^{-47}$ GeV$^4$. This value corresponds to the predicted value ($0.003$ eV)$^4$ as the dark energy on the basis of the observation of the CMB anisotropy. This result shows that it is possible to identify $a_q$ with the quintessential axion which is the origin of the dark energy.

Although the scenario is very similar to the one in [16], our model is extended by introducing additional chiral superfields. As the result, the scale of the condensations in the hidden sector can take a lower value than that in [16]. It easily makes the $\mu$ scale take an appropriate value for the electroweak symmetry breaking, keeping the supersymmetry breaking scale in the observable sector in the desirable region. Moreover, it can give an interesting explanation for the quantitative relation between the dark matter abundance and the baryon number asymmetry, as shown in the next section.

4. Abundance of baryons and dark matter

In this framework, both the abundance of the baryons and the dark matter may also be explained following the scenario presented in [9]. We discuss this point in this section.

We consider the early time evolution of the $D$-flat direction of the singlet fields $\tilde{S}_1$ and $\tilde{S}_2$ defined by

$$\langle \tilde{S}_1 \rangle = \frac{u}{\sqrt{3}} e^{i(\xi/u)}, \quad \langle \tilde{S}_2 \rangle = \frac{u}{\sqrt{2}} e^{i(\xi/u)},$$

(19)

10 In the present case we need to consider an additional supersymmetry breaking based on another source such as the modulus $F$-term breaking, since $\Lambda_h \simeq 10^{10}$ GeV is too small to realize the gaugino mass of this order through the gaugino condensation (4).

11 We do not discuss the origin of nonzero $\langle \tilde{\phi} \rangle$ but just assume its value in this paper. This value of $\langle \tilde{\phi} \rangle$ is also required from another aspect of the model as seen in section 4.
which is expected to be realized when they have effectively the same negative squared mass\(^\text{12}\). This direction is slightly lifted by the nonrenormalizable operators in the scalar potential which are induced from the first term of the superpotential \(W_{\text{mix}}\). In the following discussion, the inflaton is assumed to couple similarly with the fields in the observable and hidden sectors.

In the early universe, there are effective contributions to the scalar potential for this flat direction, which is induced by the supersymmetry breaking effects caused by the large Hubble constant \(H\) \cite{26} and the thermal effects \cite{27} other than the ordinary soft supersymmetry breaking. If we take account of these effects, the scalar potential in this direction is found to be expressed as

\[
V \simeq \left( -cH^2 + M_u^2(T) \right) u^2 + \frac{5|d_1|^2}{12} \frac{u^8}{M_{\text{pl}}^4} + \frac{1}{6\sqrt{2}} \left\{ \left( \frac{aM_{\text{susy}}e^{i\theta_a}}{M_{\text{pl}}^2} + \frac{bHe^{i\theta_b}}{M_{\text{pl}}^2} \right) u^5 e^{i(\xi/u)} + \text{h.c.} \right\},
\]

(20)

where \(M_{\text{susy}}\) is a typical soft supersymmetry breaking scale of \(O(1)\) TeV and the coefficients \(a\), \(b\) and \(c\) are \(O(1)\) real constants. CP phases \(\theta_a\) and \(\theta_b\) in the curly brackets are induced by the above mentioned supersymmetry breaking effects, which violate \(U(1)_X\) by an amount of \(\Delta X(= -4)\). The effective mass \(M_u^2(T)\) contains the usual soft supersymmetry breaking mass \(m_{\tilde{S}}^2\) of \(O(M_{\text{susy}}^2)\) and the thermal mass \(C_T \lambda_{1,2}^2 T^2\) caused by the coupling of \(\tilde{S}_{1,2}\) with \(\tilde{H}_{1,2}\) and \(\tilde{C}_{1,2}\) in the thermal plasma\(^\text{13}\). It can be expressed depending on the value of \(u\) as \cite{27}

\[
M_u^2(T) \simeq \begin{cases} 
  m_{\tilde{S}}^2 & (\lambda_{1,2} u > T), \\
  m_{\tilde{S}}^2 + C_T \lambda_{1,2}^2 T^2 & (\lambda_{1,2} u < T),
\end{cases}
\]

(21)

where \(C_T\) is a numerical factor for the thermal mass.

During the inflation, the Hubble constant contribution \(H^2\) dominates the mass of the condensate in the scalar potential (20). If the sign of this Hubble constant contribution is negative \((c > 0)\) \cite{26}, the magnitude of the condensate takes a large value such as

\[
u_1 \simeq (HM_{\text{pl}}^2)^{1/3}.
\]

(22)

On the other hand, the phase \(\xi/u\) of the condensate takes one of the five distinct values \(\xi/u = -\theta_b/5 + 2\pi \ell/5\ (\ell = 1\sim 5)\) at the potential minimum. Since \(H\) is much larger than the mass of the condensate at this period and then its evolution is almost the critical damping, the condensate follows this instantaneous potential minimum.

The dilute plasma appears as a result of a partial decay of the inflaton. Then the temperature rapidly increases to \(T_{\text{max}} \simeq (T_R^2HM_{\text{pl}})^{1/4}\) \cite{27}. \(T_R\) is the reheating

\(^{12}\)Since the potential minimum can be realized along the subspace \(3|\tilde{S}_1|^2 = 2|\tilde{S}_2|^2\) as long as the soft supersymmetry breaking masses for \(\tilde{S}_1\) and \(\tilde{S}_2\) are equal, one may expect \(|\langle \tilde{S}_1 \rangle| = u/\sqrt{3}\) and \(|\langle \tilde{S}_2 \rangle| = u/\sqrt{2}\). In the phase part of the expression in equation (19), we omit the component orthogonal to \(\xi\), which corresponds to the pNG boson associated with the \(U(1)_{X}\) breaking as discussed in section 3. It is expressed as \(\theta \equiv (-3\sqrt{2}\lambda_{1} + 2\sqrt{3}\lambda_{2})/5\). Its potential is induced not from \(W_{\text{mix}}\) but through the \(SU(3)\) and \(SU(\text{N})\) instanton effects.

\(^\text{13}\)Since the effective degrees of freedom for the relativistic fields are not same in both sectors, the reheating temperature in each sector may be somewhat different. However, this difference does not change the following discussion greatly.
temperature realized after the completion of the inflaton decay and it can be expressed as \( T_R \simeq \sqrt{M_{\text{pl}} \Gamma_1} \). \( \Gamma_1 \) is the inflaton decay width. If this temperature \( T_{\text{max}} \) does not satisfy \( \lambda_{1,2} |u_1| < T_{\text{max}} \), no thermal contribution to \( M^2_u(T) \) appears and \( M^2_u(T) \) takes the upper expression in equation (21) \([26,27]\). Thus, we can neglect the thermal effects during the inflation as long as the following lower bound on \( \lambda_{1,2} \) is satisfied:

\[
\lambda_{1,2} > T_R^{1/2} H^{-1/12} M_{\text{pl}}^{-5/12},
\]

where \( H_1 \) is the Hubble parameter during the inflation. In the following discussion, we assume that this is satisfied.

When \( H \) decreases to \( H \sim M_{\text{susy}} \) as a result of the evolution of the inflaton\(^{14} \), the effective squared mass of the condensate becomes positive and then \( u = 0 \) is the minimum of the scalar potential \( V \). The condensate starts to oscillate around \( u = 0 \), and the thermal effects due to the dilute plasma to \( M^2_u(T) \) are expected to appear. Then \( M^2_u(T) \) takes the lower expression in equation (21). At this time, the dominant term for the \( R \) evolution due to the AD mechanism \([10]\). The produced \( X \) asymmetry can be estimated by taking into account that the \( X \) current conservation is violated by the dominant \( X \) breaking operator in the curly brackets in equation (20) as

\[
\frac{d\Delta n_X(t)}{dt} = \Delta X \frac{a M_{\text{susy}} u^5}{M_{\text{pl}}} \sin \delta,
\]

where \( \Delta X \) is the \( X \) charge of that operator and \( \delta \) is determined by the difference between \( \theta_a \) and \( \theta_b \). By solving this equation, the \( X \) asymmetry produced in the condensates \( \langle \hat{S}_{1,2} \rangle \) at this period is found to be roughly expressed as \([10,26,27]\)^{15}

\[
\Delta n_X(t) \simeq \Delta X \frac{M_{\text{susy}}}{H} H^{5/3} M_{\text{pl}}^{1/3} \sin \delta,
\]

where \( t \) is the time when \( H \sim M_{\text{susy}} \). Following this period, the reheating due to the inflaton decay is completed at \( H \sim \Gamma_1 \).

Since the oscillation of the condensate behaves as the matter for the expansion of the universe, it can dominate the energy density of the universe before its decay, which is considered to occur at \( H \sim \Gamma_{\hat{S}_{1,2}} \). This is the case in the present model, since \( \Gamma_{\hat{S}_{1,2}} < \Gamma_1 \) is satisfied for reasonable values of \( \lambda_{1,2} \). The \( X \) asymmetry stored in the condensate is liberated into the thermal plasma in the observable sector and also into the hidden sector through the decay of the condensate by the \( X \) conserving couplings \( \lambda_{1,2} \hat{S}_1 H_1 H_2 \) and \( \lambda_2 \hat{S}_2 \hat{C}_1 \hat{C}_2 \), respectively. Although the \( X \) charge is assumed to be broken through the \( SU(N) \) super Yang–Mills dynamics in the hidden sector, the singlet part in the hidden sector is connected to the super-Yang–Mills part only by the gravitational interactions

\(^{14} \) Here we assume an inflation scenario such that this period is before the reheating. This means that \( M_{\text{susy}} > \Gamma_1 \) is satisfied and then \( T_R < \sqrt{M_{\text{pl}} M_{\text{susy}}} \simeq 10^{11} \text{ GeV} \).

\(^{15} \) The rigorous estimation requires the numerical calculation as discussed in [26]. It is beyond the scope of this paper and we do not refer to it further here.
as shown in $W_{\text{mix}}$. As a result, this $X$ breaking in the super-Yang–Mills part does not affect the part composed of the singlet chiral superfields where the $X$ charge is exactly conserved. The condensate decays into this $X$ conserving part. Taking account of these features, the ratio of the $X$ asymmetry $\Delta n_X^{\text{ob}}$ and $\Delta n_X^{\text{hid}}$ liberated into each sector to the entropy density $s$ is estimated as

$$Y_X^i \equiv \frac{\Delta n_X^i(\tilde{t}_R)}{s} = \frac{\Delta n_X^i(\tilde{t}_R)}{s} \approx \frac{\Delta n_X^i(t)}{2s} \frac{t^2}{\Gamma_R} \approx \frac{\Delta X}{2T_R} \frac{t^2}{\Gamma_R} \sim \frac{\Delta X}{2T_R} \frac{t^2}{\Gamma_R} \sim \frac{\Delta X}{2T_R} \frac{M_{\text{susy}}}{M_{\text{pl}}} \frac{1}{3/2} \sin \delta,$$

(26)

where we use $\tilde{t}_R \sim \Gamma_R^{-1} \sim M_{\text{pl}}/\bar{T}_R^2$. We assume that the decay widths of $\tilde{S}_{1,2}$ satisfy $\Gamma_\tilde{S}_1 \approx \Gamma_\tilde{S}_2$ to make the discussion clear. General cases will be discussed in appendix A.

Now we discuss the physical consequence of the $X$ asymmetry liberated in both sectors. At first, we focus our attention on the observable sector. If the temperature $\bar{T}_R(\approx 10^{10}\Lambda_1$ GeV) is appropriate to keep the $X$ asymmetry and $\Delta n_X^{\text{ob}} = 0$ results. However, if the $X$ and $B - L$ violating interaction in the observable sector leaves the equilibrium before the temperature reaches $T_{ss}$, the equilibrium conditions allow $\Delta n_X^{\text{ob}} \neq 0$ and a part of it is converted into the $B - L$ asymmetry. We assume this in the following. The detailed discussion about the relation between $B - L$ and $\Delta n_X^{\text{ob}}$ is presented in appendix B.

If the $B - L$ asymmetry existing at $T_{ss}$ is kept after this, the can be related to the $B$ and $L$ asymmetry as in the ordinary MSSM case as follows:

$$B = \frac{4(2N_g + 1)}{22N_g + 13} (B - L),$$

(27)

where $N_g$ is the generation number of quarks and leptons. Thus, the $B$ asymmetry produced in this scenario is finally estimated as

$$Y_B \equiv \frac{\Delta n_B}{s} \approx \frac{\Delta n_X f(N_g) \kappa}{2s} \approx \frac{\Delta X}{2T_R} \frac{M_{\text{susy}}}{M_{\text{pl}}} \frac{1}{3/2} f(N_g) \kappa \sin \delta,$$

(28)

where equation (26) is used and $\kappa \leq 1$ is introduced to take account of the washout effect. $f(N_g)$ is a numerical factor defined by

$$f(N_g) = \frac{B - L 4(2N_g + 1)}{22N_g + 13},$$

(29)

and $f(3) \approx 0.3$ for $N_g = 3$. Using this factor in equation (28), we find that this scenario can produce the currently observed $B$ asymmetry $Y_B = (0.6 - 1) \times 10^{-10}$ as long as $\bar{T}_R \gtrsim 10^4/(\Delta X \kappa \sin \delta)$ GeV is satisfied.

It is useful to comment on the relation of this scenario to the ordinary thermal leptogenesis here. In the seesaw model, the leptogenesis is usually considered on the basis of the out-of-equilibrium decay of the heavy right-handed neutrinos [30] or the decay of sneutrino condensate [31]. However, if we consider the $X$ charge asymmetry generation

16 This $\bar{T}_R$ is found to be a marginal value for the cosmological gravitino problem [28].
along the almost flat direction of $\langle \tilde{S}_1 \rangle$ as discussed above, the $B$ asymmetry produced by this usual leptogenesis might not be the dominant one. As mentioned before, we assume that the decay of the condensate is completed at the temperature above $T_{ss}$, which can be sufficiently lower than the masses of the right-handed heavy neutrinos. Then the $B$ asymmetry produced through the usual scenario seems to be washed out or overridden by the $B$ asymmetry produced in the present scenario\textsuperscript{17}.

Next we consider the hidden sector. Our interest is whether the $X$ asymmetry liberated into the hidden sector through the coupling $\lambda_2 S_2 C_1 \tilde{C}_2$ can explain the dark matter abundance. Since there is no renormalizable $U(1)_X$ violating interaction in the singlet part and also no renormalizable interaction converts the $X$ asymmetry in the singlet part into the superfields as long as no scalar components with the $X$ charge are considered to be conserved exactly within the part composed of the singlet chiral superfields as long as no scalar components with the $X$ charge get the vacuum expectation values. These features suggest that the lightest fermionic components of the singlet part can be the dark matter candidate. Since $\tilde{\phi}$ has zero $X$ charge, the $X$ charge conservation is not broken even in the case of $\langle \tilde{\phi} \rangle \neq 0$. Thus, if $\langle \tilde{\phi} \rangle \neq 0$ is realized, fermionic components of the singlet fields can get mass through the interactions in $W^S_{\text{hid}}$ without violating the $X$ charge conservation. Mass terms of these fermions can be written as

$$
(C_1 C_3) \left( \begin{array}{cc} 0 & f_1 \langle \tilde{\phi} \rangle^2 \frac{M_{pl}}{f_1} \\ f_1 \langle \tilde{\phi} \rangle^2 \frac{M_{pl}}{f_1} & 0 \end{array} \right) \left( \begin{array}{c} C_1 \\ C_3 \end{array} \right) + (C_2 C_4) \left( \begin{array}{cc} 0 & f_2 \langle \tilde{\phi} \rangle^2 \frac{M_{pl}}{f_2} \\ f_2 \langle \tilde{\phi} \rangle^2 \frac{M_{pl}}{f_2} & 0 \end{array} \right) \left( \begin{array}{c} C_2 \\ C_4 \end{array} \right)$$

$$+ (S_2 \tilde{C}_5) \left( \begin{array}{cc} 0 & f_3 \langle \tilde{\phi} \rangle^2 \frac{M_{pl}}{f_3} \\ f_3 \langle \tilde{\phi} \rangle^2 \frac{M_{pl}}{f_3} & 0 \end{array} \right) \left( \begin{array}{c} \tilde{S}_2 \\ \tilde{C}_5 \end{array} \right).$$

This shows that all fermionic components relevant to the singlet part $W^S_{\text{hid}}$ have masses $m_{\text{LP}} = O(\langle \tilde{\phi} \rangle^2 / M_{pl})$ and some of them constitute the lightest stable ones to which the $X$ charge asymmetry is finally distributed.

Taking account of these aspects and equations (26) and (28), we can estimate the ratio of the baryon energy density to this dark matter energy density in such a way as

$$\frac{\Omega_B}{\Omega_{\text{DM}}} \sim \frac{m_p Y_B}{m_{\text{LP}} Y^S_{\text{hid}}} \sim f(3) \kappa \frac{m_p}{m_{\text{LP}}},$$

where $\Omega_i$ is the ratio of the energy density $\rho_i$ to the critical energy density $\rho_{\text{cr}}$ in the present universe. Mass of the proton is represented by $m_p$. This relation suggests that the currently observed value $\Omega_B/\Omega_{\text{CDM}} \sim 0.17$ can be explained if $m_{\text{LP}}$ is a similar order value to $m_p$ in the case of $\kappa = O(1)$. This value of $m_{\text{LP}}$ can be realized for $\langle \tilde{\phi} \rangle = O(10^9-10^{10})$ GeV. It is interesting to note that this $\langle \tilde{\phi} \rangle$ is the same order as the predicted one for the explanation of the dark energy in the previous section. In the present scenario the dark matter mass is related not only to $\Omega_B/\Omega_{\text{DM}}$ but also to the amount of the dark energy. This may give a viewpoint to approach the tuning problem of the dark matter mass, which is mentioned in the footnote in section 1.

\textsuperscript{17} Even in the ordinary scenario based on the heavy right-handed (s)neutrino decay, it may be possible to make the reheating temperature low enough such as $T_R \simeq 10^6$ GeV if the neutrino mass texture satisfies suitable features [32]. In such a case both sources of the $B$ asymmetry may compete.
Finally, we give an additional comment on the $\mu$ term. As mentioned in section 3, the present model has two sources for the $\mu$ term. Here we note that the late time evolution of $\langle \tilde{S}_1 \rangle$ may give an important contribution to the $\mu$ term \cite{9}. We assume $H_1 \sim 10^{13}$ GeV during the inflation on the basis of the CMB data and $T_R \lesssim 10^9$ GeV. For these values we obtain $T_{\text{max}} \sim 10^{13}$ GeV. Thus, equation (22) gives $u_1 \sim 10^{16}$ GeV and equation (23) suggests that $\lambda_1 \gtrsim 10^{-5} T_\text{R}^{-1/2}$ should be satisfied. When the temperature decreases from $T_{\text{max}}$ to $T_e \sim m_{3/2}/\lambda_1$, $M_1^2(T)$ represented by the lower one in equation (21) starts to be dominated by the soft supersymmetry breaking mass $m_{\tilde{S}_1}^2$. If $m_{\tilde{S}_1}^2 < 0$ is realized by some reason \cite{33,34}, $\langle \tilde{S}_1 \rangle_0 \neq 0$ becomes the true vacuum after this period. Since the $\mu$ term is generated from the last term in $W_{\text{ob}}$ as $\mu = \lambda_1 \langle \tilde{S}_1 \rangle_0$, $\langle \tilde{S}_1 \rangle_0$ should be in the range $|\langle \tilde{S}_1 \rangle_0| \lesssim 10^{11} T_\text{R}^{-1/2}$ GeV to realize the appropriate $\mu$ scale for the above mentioned $\lambda_1$. Although the condensate again starts to oscillate around $\langle \tilde{S}_1 \rangle_0$, it instantaneously decays into the light fields through the $X$ conserving coupling $\tilde{S}_1 H_1 H_2$ since $H < \Gamma_{\tilde{S}_1}$ is satisfied at this time. The released energy cannot dominate the total energy density $((\pi^2/30) g_s T^4 \gg m_{3/2}^2/|\langle \tilde{S}_1 \rangle_0|^2)$ and then the effects of the produced entropy are negligible. Thus, even in this case the $X$ asymmetry obtained in equation (26) can be used as the origin of the $B$ asymmetry.

5. Summary

We have studied a possibility that all of the baryon number asymmetry, the dark matter and dark energy abundance might be related to the extended global Abelian symmetries in a supersymmetric model. The introduction of some singlet chiral superfields into the MSSM enlarges the global symmetry of the system. They may be taken as the Peccei–Quinn type symmetry $U(1)_{\text{PQ}}$ and the $R$-symmetry $U(1)_R$ or their linear combinations $U(1)_{\tilde{X}}$ and $U(1)_X$. If these global symmetries are shared in the observable and hidden sectors, they may present a unified explanation for these quantities. In this paper, such a kind of example has been concretely constructed. By using this model, we have shown that these symmetries could be a common source for them. In this scenario, the hidden sector is required to have two parts, which are connected only by the gravitational interaction. One of them is required to break the global symmetry $U(1)_{\tilde{X}}$ explicitly through the $SU(N)$ instanton effect for the dark energy explanation. Another part to prepare the massive dark matter should conserve $U(1)_X$ exactly. The flat directions associated with the introduced singlet fields play the essential role of producing the $X$ charge asymmetry.

The dark energy is explained as the quintessential axion. It is considered to appear in the model as the mixture state of the pNG boson associating with the spontaneous symmetry breaking of the global symmetry $U(1)_{\tilde{X}}$ and the model independent axion given by the superstring. The intermediate scale introduced relating to the $U(1)_{\tilde{X}}$ breaking can successfully generate the dark energy scale. It is also suitable for the explanation of the strong CP problem due to the invisible axion. In our case this invisible axion cannot be a dark matter candidate. The operators which play the important role for this explanation also produce the $\mu$ term with the suitable scale.

\footnote{Such a value of $\langle \tilde{S}_1 \rangle_0$ may be expected to be determined either by the nonrenormalizable terms or by the pure radiative symmetry breaking effect as discussed \cite{34}.}
Both the baryons and the dark matter are considered to be produced from the charge asymmetry of $U(1)_X$. This asymmetry is stored in the condensate of the singlet scalar components due to the AD mechanism. A part of this asymmetry is converted into the $B$ asymmetry in the observable sector. Another part is liberated into the hidden sector and causes the dark matter abundance. We have shown that both observed values of $Y_B$ and $\Omega_B/\Omega_{CDM}$ can be explained by them in a consistent way. However, it may be fair to say that our scenario does not evade the problem mentioned in footnote 1. Nevertheless, we should note that it can give the relation between the mass of the dark matter and the number densities of the dark matter and the baryon on the basis of the dynamics of the common flat directions related to the extended global symmetry. We think that this kind of approach may give a direction to consider these problems.

Since the dark matter in this model interacts with the fields in the observable sector only through gravity, it can be distinguished from the ordinary dark matter candidates such as the neutralino and the axion. However, its detection is difficult. The experimental signature of the model might be found in the phenomena related to the coupling $\lambda_1 \tilde{S}_1 \tilde{H}_1 \tilde{H}_2$, that is, the phenomena including the neutral Higgs scalar or the neutralinos. To proceed with such a study, it is necessary to fix the model in a more detailed way.

In this paper we assume a rather simple structure for the hidden sector, in particular, for the part required for the dark matter mass generation. In the present framework, however, it may be possible and also natural to suppose a similar structure to the observable sector for the hidden sector except for the $SU(N)$ super-Yang–Mills part. This may constitute the mirror world for the observable sector. In this case the dark matter abundance may be explained in just the same way as the $B$ asymmetry in this scenario. The study of this kind of possibility may be worth proceeding with.

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**Appendix A**

In this appendix we discuss the amount of $X$ asymmetry liberated into each sector for the general cases with arbitrary $\Gamma_{\tilde{S}_1}$ and $\Gamma_{\tilde{S}_2}$. If we take into account that the decay products of the condensate behave as radiation and also no photon is produced through its decay into the hidden sector, we can estimate the $X$ asymmetry $\Delta n_X^\text{ob}(t_R)$ generated in the observable sector in the case $\Gamma_{\tilde{S}_1} > \Gamma_{\tilde{S}_2}$ as

$$\frac{\Delta n_X^\text{ob}(t_R)}{s} = \frac{\Delta n_X(t)}{s} \frac{\Gamma_{\tilde{S}_1}}{\Gamma_{\tilde{S}_1} + \Gamma_{\tilde{S}_2}} \left( \frac{t}{t_1} \right)^2,$$

and also in the case $\Gamma_{\tilde{S}_1} < \Gamma_{\tilde{S}_2}$ as

$$\frac{\Delta n_X^\text{ob}(t_R)}{s} = \frac{\Delta n_X(t)}{s} \frac{\Gamma_{\tilde{S}_1}}{\Gamma_{\tilde{S}_1} + \Gamma_{\tilde{S}_2}} \left( \frac{t_1}{t_2} \right)^2 \left( \frac{t_2}{t_1} \right)^{3/2},$$

(A.2)
where $\tilde{t}_R \simeq t_1 = \Gamma^{-1}_{\tilde{s}_1}$ and $t_2 = \Gamma^{-1}_{\tilde{s}_2}$. Although $Y_B$ has a similar expression to that discussed in the text for the case of $\Gamma_{\tilde{s}_1} > \Gamma_{\tilde{s}_2}$, there appears an additional suppression factor $(\frac{\Gamma_{\tilde{s}_1}}{\Gamma_{\tilde{s}_2}})^{1/2}$ for the case of $\Gamma_{\tilde{s}_1} < \Gamma_{\tilde{s}_2}$.

$$Y_B \simeq \Delta X \tilde{T}_R \frac{1}{M_{\text{pl}}^{2/3}} \left( \frac{\Gamma_{\tilde{s}_1}}{\Gamma_{\tilde{s}_2}} \right)^{1/2} f\left( N_{g} \right) \kappa \sin \delta.$$  (A.3)

The higher reheating temperature $\tilde{T}_R$ is required to produce the same $Y_B$ in comparison with the case of $\Gamma_{\tilde{s}_1} \gg \Gamma_{\tilde{s}_2}$.

The ratio of the energy density of the baryons and the dark matter can be expressed in both cases as

$$\frac{\Omega_B}{\Omega_{\text{DM}}} \simeq \frac{m_p Y_B}{m_{\text{LP}} Y_{\text{hid}}} \simeq f\left(3\right) \frac{m_p}{m_{\text{LP}}} \frac{\Gamma_{\tilde{s}_1}}{\Gamma_{\tilde{s}_2}}.$$  (A.4)

This suggests that the currently observed value of $\Omega_B/\Omega_{\text{DM}}$ seems to be explained as long as $m_{\text{LP}} \simeq (\frac{\Gamma_{\tilde{s}_1}}{\Gamma_{\tilde{s}_2}})^{1/2} \kappa$ GeV is satisfied. Thus, in the present model $\Gamma_{\tilde{s}_1} > \Gamma_{\tilde{s}_2}$ is favoured from the viewpoint of the explanation of the dark energy.

### Appendix B

In this appendix we discuss how the $X$ asymmetry in the observable sector is connected to the $B$ asymmetry [9]. For this study, we use the detailed balance equations of relevant interactions [35,36]. The particle–antiparticle number asymmetry $\Delta n_f$ in the case of $\mu_f \ll T$ can be approximately represented as

$$\Delta n_f \equiv n_f - n_{f^c} = \begin{cases} \frac{g_f T^2 \mu_f}{6} & (f: \text{fermion}), \\ \frac{g_f T^2 \mu_f}{3} & (f: \text{boson}), \end{cases}$$  (B.1)

where $\mu_f$ is chemical potential and $g_f$ is a number of relevant internal degrees of freedom of the field $f$. By solving the detailed valence equations for the chemical potential $\mu_f$, we can estimate the charge asymmetry.

If the $SU(2)$ and $SU(3)$ sphaleron interactions are in thermal equilibrium, we have conditions such as

$$\sum_{i=1}^{N_g} (3\mu_{Q_i} + \mu_{L_i}) + \mu_{H_1} + \mu_{H_2} + 4\mu_W = 0, \quad \sum_{i=1}^{N_g} (2\mu_{Q_i} - \mu_{U_i} - \mu_{D_i}) + 6\mu_{\tilde{\beta}} = 0,$$  (B.2)

where $N_g$ is a number of the generation of quarks and leptons. The cancellation of the total hypercharge or the electric charge of plasma in the universe requires

$$\sum_{i=1}^{N_g} (\mu_{Q_i} + 2\mu_{U_i} - \mu_{D_i} - \mu_{L_i} - \mu_{E_i}) + \mu_{H_2} - \mu_{H_1}$$

$$+ 2 \sum_{i=1}^{N} (\mu_{Q_i} + 2\mu_{U_i} - \mu_{D_i} - \mu_{E_i}) + 2 (\mu_{H_2} - \mu_{H_1}) = 0.$$  (B.3)
When Yukawa interactions in $W_{ab}$ are in thermal equilibrium, they impose the following conditions\(^\text{19}\):  
\[
\begin{align*}
\mu_{Q_i} - \mu_{U_j} + \mu_{H_2} &= 0, \\
\mu_{L_i} - \mu_{E_j} + \mu_{H_1} &= 0, \\
\mu_{S_i} + \mu_{H_1} + \mu_{H_2} &= 0.
\end{align*}
\]  
(B.4)
There are also the conditions for the gauge interactions in the thermal equilibrium, which are summarized as  
\[
\mu_{\tilde{Q}_i} = \mu_{\tilde{\beta}} + \mu_{Q_i} = \mu_{\tilde{W}} + \mu_{Q_i} = \mu_{\tilde{B}} + \mu_{Q_i},
\]  
(B.5)
where $\mu_{\tilde{\beta}}$, $\mu_{\tilde{W}}$ and $\mu_{\tilde{B}}$ stand for gauginos in the MSSM. Similar relations to equation (B.5) are satisfied for leptons $\tilde{L}_i$, doublet Higgs fields $\tilde{H}_{1,2}$ and other fields $\tilde{U}_i, \tilde{D}_i, \tilde{E}_i$ which have the SM gauge interactions. Flavour mixings of quarks and leptons due to the Yukawa couplings allow us to consider the flavour independent chemical potential such as $\mu_{\tilde{Q}} = \mu_{Q_i}$ and $\mu_L = \mu_{L_i}$.

If an operator violating both $B - L$ and $X$ exists, only a linear combination of these two $U(1)$s is absolutely conserved. Then a part of the $X$ asymmetry can be converted into the $B - L$ asymmetry. We consider an effective operator $(y^2/M_R)(\tilde{L}\tilde{H}_2)^2$ for such an example\(^\text{20}\). The thermal equilibrium condition of this operator can be written as  
\[
\mu_L + \mu_{H_2} = 0.
\]  
(B.6)
We have to assume that the effective operator $(y^2/M_R)(\tilde{L}\tilde{H}_2)^2$ is in thermal equilibrium at first and then leaves the thermal equilibrium before the time when the soft supersymmetry breaking operators start being in thermal equilibrium as discussed in the text. Since the reaction rates for this operator and the soft supersymmetry breaking operators are estimated as $\Gamma_\nu \sim y^4T^3/M_R^2$ and $\Gamma_{ss} \sim m_{3/2}^2/T$ \(^\text{17}\), this assumption requires that $\Gamma_\nu > H$ should be satisfied for the temperature $T_\nu$ which satisfies the condition $H \gtrsim \Gamma_{ss}$. The temperature $T_\nu$ can be estimated as  
\[
T_\nu \gtrsim \frac{10\langle H_2 \rangle^4}{m_\nu^2M_{pl}} \sim 10^9 \left( \frac{\text{eV}}{m_\nu} \right)^2 \sin^4 \beta \text{ GeV},
\]  
(B.7)
where $m_\nu$ is the neutrino mass and $\tan \beta = |\langle \tilde{H}_2 \rangle/\langle \tilde{H}_1 \rangle|$. Since the temperature $T_\nu$ satisfies $T_\nu \gtrsim T_{ss} \simeq 10^7 \text{ GeV}$,\(^\text{21}\) we can have an independent chemical potential from the thermal equilibrium conditions (B.2)–(B.6). It can be taken as $\mu_{H_2}$, which can be related to the above mentioned remaining symmetry. On the other hand, the reheating temperature $T_R$ should satisfy $T_R > T_\nu$. This means that the present scenario can have as much trouble with the gravitino bound as the standard leptogenesis if we take account of the cosmological neutrino mass bound \(^\text{1}\) other than the conditions required by the neutrino oscillation data \(^\text{37, 38}\).

\(^\text{19}\) We should note that each term in $W_{\text{mix}}$ leaves the thermal equilibrium at $T \sim M_{\text{pl}}$. Since $\tilde{S}_1$ has no other coupling to the MSSM contents than $\lambda_1 \tilde{S}_1 \tilde{H}_1 \tilde{H}_2$, the last one in equation (B.4) is the only condition for $\mu_{\tilde{S}_1}$.

\(^\text{20}\) It corresponds to the effective neutrino mass operator induced in the ordinary seesaw mechanism. It is obtained from the operator $yN\tilde{L}\tilde{H}_2$ by integrating out the heavy right-handed neutrino $\tilde{N}$ with the mass $M_R$, where Yukawa coupling constants are represented by $y$ and the generation indices are abbreviated.

\(^\text{21}\) For $T_\nu \lesssim T_{ss}$, $\mu_\beta = 0$ is satisfied and then equations (B.2)–(B.6) result in $\mu_{H_2} = 0$. The $X$ asymmetry produced in the observable sector through the decay of the condensate disappears.
By solving equations (B.2)–(B.6), $\mu_Q$, $\mu_L$, $\mu_{H_1,2}$ and $\mu_B$ can be written with the chemical potential of Higgsino field $\tilde{H}_2$ at temperature $T$ ($\gtrsim T_{ss}$) in such a way as

$$\mu_Q = \frac{17N_g}{N_g(10N_g^2 - 17N_g - 15)}H_2,$$

$$\mu_L = -\frac{5(4N_g + 3)}{10N_g^2 - 17N_g - 15}\mu_H_2,$$

$$\mu_{H_1} = -\frac{10N_g + 3}{10N_g^2 - 17N_g - 15}H_2,$$

$$\mu_B = -\frac{(10N_g + 3)N_g}{10N_g^2 - 17N_g - 15}H_2.$$  

By defining $B$ and $L$ as $\Delta n_B \equiv BT^2/6$ and $\Delta n_L \equiv LT^2/6$, we can calculate these values at $T_{ss}$ by using equations (B.1) and (B.8). Using the results, we can relate the $X$ asymmetry in the observable sector to the $B - L$ asymmetry as follows:

$$B - L = \frac{20N_g^3 + 162N_g^2 - 24N_g - 72}{360N_g^3 + 3308N_g^2 - 11419N_g - 1413}X,$$  

where $X$ stands for the $X$ asymmetry liberated into the observable sector which is defined by $\Delta n_X^\text{ob} \equiv XT^2/6$. After the freeze-out of $(LH_2)^2$ at the temperature $T$ ($\gtrsim T_{ss}$), the $B - L$ asymmetry is kept conserved.

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