Inflationary baryogenesis with low reheating temperature and testable neutrino-antineutron oscillation

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Recently we extended the standard model by four TeV-scale fields including a singlet fermion, an isotriplet and two isosinglet diquark scalars to generate the cosmological baryon asymmetry with an observable neutron-antineutron oscillation. We now supersymmetrize our model but do not constrain it at the TeV scale. The superpartner of the singlet fermion can serve as an inflaton field. Its three-body decays, mediated by the isosinglet diquarks and their superpartners, can simultaneously provide a low reheating temperature and a sizable CP asymmetry. We thus can realize a nonthermal baryogenesis without the gravitino problem. Meanwhile, we can have a testable neutron-antineutron oscillation induced by the exchange of one isosinglet and two isotriplet diquarks if the isodiquark diquark is at the TeV scale.

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Many baryogenesis scenarios\textsuperscript{1-7} have been studied to explain the observed baryon asymmetry in the present universe. A successful baryogenesis model should respect the Sakharov conditions\textsuperscript{8} (baryon number nonconservation, C and CP violation as well as departure from equilibrium) unless CPT is not invariant\textsuperscript{4}. If the baryogenesis works above the weak scale, the baryon number violation should be modified to be a baryon-minus-lepton number violation in the presence of the sphalerons\textsuperscript{9}, which violate the baryon and lepton numbers by an equal amount. The sphaleron processes will not affect any primordial $B-L$ asymmetry and will convert the $B-L$ asymmetry to a baryon asymmetry and a lepton asymmetry. The $B-L$ asymmetry can be composed of a pure baryon asymmetry or a pure lepton asymmetry or any unequal baryon and lepton asymmetries. The baryon and/or lepton number violation can lead to other interesting phenomenologies. For example, the baryon number violation of two units could result in a neutron-antineutron oscillation\textsuperscript{10}, the lepton number violation of two units could result in a neutrinoless double beta decay\textsuperscript{11}, whereas the combined baryon and lepton number violation could result in a proton decay\textsuperscript{12}.

In a recent work\textsuperscript{13}, we extended the $SU(3)_c \times SU(2)_L \times U(1)_Y$ standard model (SM) by four TeV-scale fields (a singlet fermion, an isotriplet and two isosinglet diquark scalars) to generate the cosmological baryon asymmetry with an observable neutron-antineutron oscillation. In this paper we will study the supersymmetric version of our model but will not constrain it at the TeV scale. Like the sneutrinos (the superpartner of the right-handed neutrinos in the supersymmetric seesaw\textsuperscript{14} model), the superpartner of the singlet fermion can drive a chaotic inflation\textsuperscript{15}. Through the three-body decays\textsuperscript{16} of the inflaton, the universe can be reheated at a low temperature to avoid the gravitino problem\textsuperscript{17}. At the same time, we can obtain a sizable CP asymmetry in the inflaton decays to realize a nonthermal baryogenesis. Furthermore, we can have a neutron-antineutron oscillation by the exchange of one isosinglet and two isotriplet diquarks. This neutron-antineutron oscillation can be sensitive to the future experiments if the isotriplet diquark is at the TeV scale.

For simplicity, we only write down the superpotential relevant to our demonstration,

\begin{equation}
W \supset y_{ai}\hat{\Delta}_a \hat{D}_i^c \hat{X}^c + \bar{f}_{aij} \hat{\Delta}_a \hat{D}_i \hat{D}_j + h_{ai} \hat{Q}_i \tau_2 \hat{\Omega}_j
+ \kappa_{ij} \hat{\Delta}_a \text{Tr}(\hat{\Omega}\hat{\Omega}) + \bar{\kappa}_a \hat{\Delta}_a \text{Tr}(\hat{\Omega}\hat{\Omega}) + \frac{1}{2} M_X \hat{X}^c \hat{X}^c
+ M_a \hat{\Delta}_a + M_\Omega \text{Tr}(\hat{\Omega}\hat{\Omega}) .
\end{equation}

Here

\begin{equation}
\hat{X}(1, 1, 0)
\end{equation}

is a singlet superfield,

\begin{equation}
\hat{\Delta}(3, 1, +\frac{2}{3}), \hat{\Omega}(3, 3, -\frac{1}{3}) = \begin{bmatrix}
-\frac{1}{\sqrt{2}} \hat{\omega}_{\frac{3}{2}} & -\frac{1}{2} \hat{\omega}_{\frac{1}{2}} \\
\frac{1}{\sqrt{2}} \hat{\omega}_{\frac{3}{2}} & \hat{\omega}_{\frac{1}{2}}
\end{bmatrix},
\end{equation}

\begin{equation}
\hat{\Delta}(3, 1, -\frac{2}{3}), \hat{\Omega}(3, 3, +\frac{1}{3}) = \begin{bmatrix}
-\frac{1}{\sqrt{2}} \hat{\omega}_{\frac{3}{2}} & \frac{1}{2} \hat{\omega}_{\frac{1}{2}} \\
\frac{1}{\sqrt{2}} \hat{\omega}_{\frac{3}{2}} & \hat{\omega}_{\frac{1}{2}}
\end{bmatrix}
\end{equation}

are the diquark superfields, while

\begin{equation}
\hat{Q}(3, 2, 0), \hat{U}(3, 1, 2), \hat{D}(3, 1, -\frac{1}{3})
\end{equation}

denote the usual quark superfields. Corresponding to the baryon number $B = \frac{1}{3}$ of the quark superfields, we
assign the baryon number $B = 1$ for the singlet superfield while the baryon number $B = -1$ for the diquark superfields ($\hat{A}, \hat{A}^c$) and ($\hat{Q}, \hat{Q}^c$). The singlet superfield is forbidden to have the couplings with the lepton and Higgs superfields so that we can avoid the dangerous proton decay. For this purpose, we can introduce certain discrete symmetries, such as a $\mathbb{Z}_2$ symmetry under which the lepton superfields (including the right-handed neutrino superfields for the seesaw mechanism) are odd while the others are even. The superpotential will yield a lagrangian as below,

$$\mathcal{L} \supset -y_{ai}(\delta_{a} u_{Ri} X^c_{R} + \bar{X}^a u_{Ri} \tilde{\delta}_{La} + \bar{u}^*_a \tilde{X}_R \delta_{La}$$

$$\quad + M_{a} \delta^*_{La} \bar{X}_R) - f_{aij}(\delta_{a} d_{Ri} \bar{d}^*_R + 2d_{Ri} \delta_{La} \bar{d}^*_R$$

$$\quad + M_{a} \delta^*_{La} d_{Ri} \bar{d}^*_R) + h_{ij}(\bar{q}_{Li} \tau_2 w q_{Lj} + 2\bar{q}_{Li} \tau_2 \bar{w} \tilde{q}_{Lj}$$

$$\quad + M_{ij} \bar{q}_{Li} \tau_2 \bar{w}) - \kappa_{a}[\delta_{a} \bar{\omega} \omega L]$$

$$\quad + 2 \text{Tr}(\bar{\omega} \omega \delta_{La} + \delta^* a \text{Tr}(\omega^2) - \bar{\omega} a \text{Tr}(\bar{\omega} \omega L)$$

$$\quad + 2 \text{Tr}(\bar{\omega} \omega \delta_{La} + \delta^* a \text{Tr}(\omega^2) - \frac{1}{2} M_X \bar{X}_R \bar{X}_R^c$$

$$\quad - \frac{\delta^*_{La} \bar{X}_R}{M_X}, - M_{ij} \text{Tr}(\bar{\omega} \omega L) + \text{H.c.} - M_{a} \tilde{X}^* \bar{X}$$

$$\quad - M_{a} \delta^*_{La} + \delta a^* \delta a) - M^2 \text{Tr}(\omega^4)$$

$$\quad + \text{Tr}(\omega^3 \omega)] \right). \tag{5}$$

The singlet scalar $\tilde{X}$ can drive a chaotic inflation, like the sneutrino in the supersymmetric seesaw model \cite{11}. Specifically, the inflaton $\tilde{X}$ will begin to oscillate at the time $t \sim 1/H(T) \sim 1/M_X$. Here the Hubble constant $H$ is given by

$$H(T) = \left(\frac{8\pi^3 g_*}{90}\right)^{1/2} \frac{T^2}{M_{Pl}} \tag{6}$$

with $M_{Pl} = \mathcal{O}(10^{19} \text{GeV})$ being the Planck mass and $g_* = \mathcal{O}(200)$ being the relativistic degrees of freedom. Subsequently, the inflaton will start to decay at the time $t \sim 1/H(T) \sim 1/\Gamma$ with $\Gamma$ being the decay width. The universe then can be reheated by the relativistic decay products. The reheating temperature is determined by

$$\Gamma = H(T) \Rightarrow T_R = \left(\frac{90}{8\pi^3 g_*}\right)^{1/4} M_{Pl} \Gamma. \tag{7}$$

The masses in the superpotential \cite{11} and the lagrangian \cite{5} are assumed to hold the following hierarchy,

$$2M_\Omega < M_X < M_a. \tag{8}$$

This means that the inflaton can only have the three-body decays mediated by the isosinglet diquark scalars and their superpartners. We show the inflaton decays at tree level and one-loop order in Fig. 1. The final states of the inflaton decays carry a baryon number $B = +1$ or $B = -1$, i.e.

$$B = +1 \quad B = -1$$

$$\tilde{X}^* \rightarrow u_R d_R \bar{d}^*_R, \quad \tilde{X} \rightarrow u_R d_R \bar{d}^*_R,$$

$$\tilde{X} \rightarrow u_R \bar{d}^*_L \omega^*, \quad \tilde{X}^* \rightarrow u_R \bar{d}^*_L \omega^*, \tag{9}$$

$$\tilde{X} \rightarrow u_R \bar{d}^*_L \omega^*, \quad \tilde{X}^* \rightarrow u_R \bar{d}^*_L \omega^*, \quad \tilde{X} \rightarrow u_R \bar{d}^*_L \omega^*,$$

$$\tilde{X} \rightarrow u_R \bar{d}^*_L \omega^*, \quad \tilde{X}^* \rightarrow u_R \bar{d}^*_L \omega^*, \quad \tilde{X} \rightarrow u_R \bar{d}^*_L \omega^*.$$

The $B = +1$ processes will generate a positive baryon number while the $B = -1$ processes will generate a negative baryon number. Therefore, the $B = +1$ processes and the $B = -1$ processes can have different decay widths to induce a net baryon number if CP is not conserved. In addition to the three-body decays, there are $3 \leftrightarrow 3$ and $2 \leftrightarrow 2$ scattering processes violating the baryon number. As long as the reheating temperature is far below the inflaton mass, the scattering processes will not wash out the induced baryon asymmetry since they have completely decoupled before the inflaton decays. We further assume that other baryon or lepton number violating interactions does not exist or have already decoupled. The baryon asymmetry from the inflaton decays then should be \cite{11},

$$\frac{n_B}{s} = \frac{\hat{\xi} n_{\tilde{X}}}{s} \bigg|_{T = T_R} \simeq \frac{3}{4} \frac{T_B}{M_X}. \tag{10}$$

Through the sphalerons, we can obtain a final baryon asymmetry \cite{11},

$$\eta_B = \frac{8}{23} \frac{n_B}{s}. \tag{11}$$

The final baryon asymmetry is determined by the reheating temperature and the CP asymmetry for a given inflaton mass. We hence calculate the decay width at tree level and the CP asymmetry at one-loop order,
FIG. 1: The three-body decays of the inflaton at tree level and one-loop order. The CP conjugation is not shown for simplicity.
FIG. 2: The six-quark interactions violating the baryon number of two units. The CP conjugation is not shown for simplicity.

\[ \tilde{\Gamma} = \Gamma X = \Gamma X \to \bar{u}_R^c d_R^c + \Gamma X \to \bar{u}_L^c \bar{d}_L^c + \Gamma X \to \bar{u}_R^c \bar{d}_R^c + \Gamma X \to u_R^c \bar{d}_R^c \]  
\[ = \frac{3}{2π} \frac{\sum ab}{\sum k} \left( \frac{\sum ab}{\sum k} \right) \Gamma \sum_{\mu} \tilde{\Gamma}_{\mu} \left( 2 - \kappa_{\mu}^a \kappa_{\mu}^b \right) r_a r_b + \kappa_{\mu}^a \kappa_{\mu}^b \left( 6 + r_a^2 r_b^2 \right) \]  
\[ \tilde{\varepsilon} = \frac{1}{2π} \frac{\sum abc}{\sum k} \left( \frac{\sum ab}{\sum k} \right) \Gamma \sum_{\mu} \tilde{\varepsilon}_{\mu} \left( 2 - \kappa_{\mu}^a \kappa_{\mu}^b \right) r_a r_b + \kappa_{\mu}^a \kappa_{\mu}^b \left( 6 + r_a^2 r_b^2 \right) \]  

where the parameter and the CP asymmetry by

\[ r_a = \frac{M_X}{M_a} \]  
\[ \tilde{\Gamma} = \frac{3}{2π} AM_X, \quad \tilde{\varepsilon} = \frac{1}{2π} \frac{B}{A}, \]  

has been defined. We further specify the decay width where the quantities
\[ A = \sum_k y_{1k}^2 \left( 2 \left( \sum_{ij} \tilde{f}_{1ij}^2 + 3\kappa_1^2 \right) r_1^2 + \kappa_1^2 \left( 6 + r_1^2 \right) \right) + \sum_k y_{2k}^2 \left( 2 \left( \sum_{ij} \tilde{f}_{2ij}^2 + 3\kappa_2^2 \right) r_2^2 + \kappa_2^2 \left( 6 + r_2^2 \right) \right) \]

\[ + 2 \sum_k y_{1k} y_{2k} \left\{ 2 \left[ \sum_{ij} \tilde{f}_{1ij} \tilde{f}_{2ij} \cos (\alpha_k + \beta_{ij}) + 3\kappa_1'\kappa_2' \cos (\alpha_k + \gamma) \right] r_1 r_2 + \kappa_1'\kappa_2' \cos (\alpha_k - \gamma) \left( 6 + r_1^2 r_2^2 \right) \right\} \]

\[ B = \sum_k y_{1k} y_{2k} \kappa_1' \kappa_2' \left[ \left( \sum_{ij} \tilde{f}_{1ij}^2 + 3\kappa_1^2 \right) (2 - r_2^2) r_1^2 - (\tilde{f}_{2ij}^2 + 3\kappa_2^2) (2 - r_1^2) r_2^2 \right] \sin (\alpha_k - \gamma) \]

\[ + \sum_k y_{1k} y_{2k} \left[ \sum_{ij} \tilde{f}_{1ij} \tilde{f}_{2ij} \sin (\alpha_k + \beta_{ij}) + 3\kappa_1'\kappa_2' \sin (\alpha_k + \gamma) \right] \left[ \kappa_1^2 (2 - r_2^2) - \kappa_2^2 (2 - r_1^2) \right] r_1 r_2 \]

\[ + \kappa_1' \kappa_2' \left[ \sum_{ij} \tilde{f}_{1ij} \tilde{f}_{2ij} \sin (\beta_{ij} + \gamma) + 3\kappa_1'\kappa_2' \sin (\gamma + \gamma) \right] \left\{ \sum_k [y_{1k}^2 (2 - r_1^2) - y_{2k}^2 (2 - r_2^2)] \right\} r_1 r_2, \]

are determined by the parameters,

\[ y_{ak} = y_{ak} e^{i\alpha_{ak}}, \quad \alpha_k = \alpha_{1k} - \alpha_{2k}, \]

\[ \tilde{f}_{aij} = \tilde{f}_{aij} e^{i\beta_{aij}}, \quad \beta_{ij} = \beta_{1ij} - \beta_{2ij}, \]

\[ \kappa_1 = \kappa_1' e^{i\gamma_1}, \quad \gamma = \gamma_1 - \gamma_2, \]

\[ \kappa_2 = \kappa_2' e^{i\gamma_2}. \]

We now indicate that our model can simultaneously provide a desired baryon asymmetry and a low reheating temperature. For example, we take

\[ y_{1k} = y_{2k} = y', \quad \tilde{f}_{1ij} = \tilde{f}_{2ij} = \tilde{f}', \]

\[ \kappa_1' = \frac{1}{2} \kappa_2' = \kappa', \quad \kappa_1 = \kappa_2 = \kappa', \]

\[ \alpha_k = \beta_{ij} = \gamma = \frac{1}{2} \delta, \quad r_1 = r_2 = r, \]

(18)

(19)

to derive

\[ T_R = \frac{9}{32\pi^2} \left( \frac{5}{\pi g_4} \right)^\frac{1}{4} y' \left\{ 4 (3\tilde{f}'^2 + \kappa'^2) r^2 (1 + \cos \delta) + 3\kappa'^2 \left( 6 + r^4 \right) \right\} \]

\[ \left[ M_X M_{PL} \right]^\frac{1}{4}, \]

(20)

\[ \tilde{\varepsilon}_X = \frac{3\kappa'^2 (3\tilde{f}'^2 + \kappa'^2) (2 - r^2) r^2 \sin \delta}{2\pi^4 \left( 3\tilde{f}'^2 + \kappa'^2 \right) r^2 (1 + \cos \delta) + 3\kappa'^2 \left( 6 + r^4 \right)}. \]

(21)

It is easy to understand that the parameter \( y' \) in the reheating temperature does not appear in the CP asymmetry since the CP asymmetry is from the interference between the tree and one-loop diagrams of the three-body decays. So, we can obtain a sizable CP asymmetry even if the \( y' \) is smaller to lower the reheating temperature.
With the previous parameter choice for the inflationary baryogenesis, the induced neutron-antineutron oscillation can arrive at a testable level \([10]\) as it has a strength of the order of
\[
G_{n\bar{n}} \sim \sum_a \frac{6\kappa_a f_{h_{11}^a} h_{11}^a}{M_a M_{h_{11}}^4} \sim \frac{12\kappa' f^2 h_{11}^2}{M_X M_{h_{11}}^4} \approx 3 \times 10^{-28}\ \text{GeV}^{-5} \left( \frac{r}{0.1} \left( \frac{\kappa'}{0.5} \right) \left( \frac{f}{0.5} \right) \right) \times \left( \frac{h_{11}}{0.1} \right)^2 \left( \frac{M_X}{10^{13}\ \text{GeV}} \right)^{-1} \left( \frac{M_{h_{11}}}{1 \text{TeV}} \right)^{-4}
\]

In this paper we studied a supersymmetric model with the singlet and diquark superfields. The scalar component of the singlet superfield can drive a chaotic inflation. The inflaton only have the three-body decays mediated by the isosinglet diquark scalars and their superpartners. The inflaton decays can simultaneously allow a low reheating temperature to avoid the gravitino problem and a sizable CP asymmetry to generate the cosmological baryon asymmetry. There will emerge an observable neutron-antineutron oscillation if the isotriplet diquark scalar is at the TeV scale. The isotriplet diquark could be verified at colliders such as the LHC \([20]\). Furthermore, our model can be modified and extended. For example, the isotriplet diquark superfields can be replaced by the isosinglet ones. We can also introduce more singlet superfields to give other inflationary scenarios, like the double sneutrino inflation \([21]\) and the hybrid sneutrino inflation \([22]\).

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