Abstract

The large tensor-to-scalar ratio reported by BICEP2 collaboration may lead to distinctive phenomenology of high-energy scale. Assuming the same origin of SUSY breaking between inflation and MSSM, we show model independent features in such high-scale SUSY. The simplest hybrid inflation, together with a new linear term for inflaton field which is induced by large gravitino mass, is excluded by BICEP2 data. For superpartner masses far above electroweak scale we estimate the reheating temperature $T_R$ after inflation. We find that $T_R$ might be beneath the value required by thermal leptogenesis if inflaton decays to its products perturbatively, but above it if non-perturbatively instead. Due to kinematically blocking effect the gravitino overproduction can be also evaded in high-scale SUSY.
1 Introduction

After the discovery\cite{1} of standard model (SM) Higgs boson at the Large Hadron collider (LHC),
low-scale supersymmetry (SUSY)\cite{2} which is favored by naturalness\cite{3} has been extensively
explored. These expositions show its difficulties in both theoretic explanation of 125 GeV Higgs
mass and experimental fits to the LHC data. Waiting for the second run of LHC is the best option
to address a possible answer on low-scale SUSY. Even so some efforts have been devoted to study
high-scale SUSY instead.

The image that high-scale SUSY deserves our attention is further emphasized by BICEP2 data.
The BICEP2 collaboration is devoted to measure Cosmic Microwave Background (CMB) temper-
ature anisotropy and polarization, which is useful to uncover the nature of inflation in the early
universe. BICEP2 collaboration reported a large tensor-to-scalar ratio $r$ of order $0.1 \sim 0.2$\cite{4},
which is against previous results\cite{5,6} shown by Plank collaboration. Since the inflation happens
at rather high energy scale, it can be applied to the study of high-scale SUSY which is beyond the
scope of colliders such as LHC.

It is possible that the signal observed by BICEP2 attributes to the dust, see, e.g., discussions in
\cite{7} and recent Plank report\cite{8} on this issue. However, if $r$ isn’t reduced much from the BICEP2 re-
ported value $\sim 0.2$\cite{1} meaningful results about high-scale SUSY can be still derived. As for model
building of inflation, a few works\cite{9,10,11,12} have shown that BICEP data favors quadratic
inflation. For reviews on inflation including quadratic inflation, see, e.g.,\cite{13,14}. Among other
things a few model independent results can be also obtained, which include high SUSY breaking
scale of order $\sim 10^{16}$ GeV, large inflaton mass $m_{\phi}$ of order $\sim 10^{13}$ GeV.

In this paper, we explore high-scale SUSY based on the assumption\cite{9,10} that inflation and
visible (namely minimal supersymmetric standard model (MSSM)) share the same origin of SUSY
breaking. In other words, the scales of SUSY breaking between these two sectors are the same.
This assumption is rational, as it can be realized in model building. Moreover, it allows us to
discuss reheating in the early universe after inflation, once the superpartner spectrum and inflaton
decay are identified explicitly.

The paper is organized as follows. In section 2 we re-analyze the model independent conse-
quences from above assumption. If one adopts a constant superpotential to cancel out the vacuum

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\footnote{Since the energy scale of inflation is proportional to $r^{1/4}$, its suppression isn’t sensitive to the change of $r$. For example, the energy scale is suppressed only by one order of magnitude for $r \sim 10^{-5}$ instead of $r \sim 0.1$. So, the motivations for high-scale SUSY remain. This note is the first step towards to study of early universe with high-scale SUSY.}
energy density so as to explain the smallness of cosmological constant (cc), the gravitino mass is found to be of order \( \sim 6m_\phi \).

In section 3, we consider hybrid inflation as an example in the course of high-scale SUSY breaking. We will show that a new linear term for inflaton field with a large coefficient proportional to \( m_{3/2} \) affects the inflation significantly, and the simplest hybrid inflation is impossible to explain all experimental data simultaneously.

In section 4, we discuss the reheating in early universe after inflation. In particular, reheating temperature \( T_R \) after inflation is estimated for superpartner mass spectrum \( m_0 > \mathcal{O}(100) \) TeV. We find that \( T_R \) might be beneath the value \( \sim 10^9 \) GeV required by thermal leptogenesis if inflaton decays to its products perturbatively, but above it if non-perturbatively instead. The gravitino overproduction in conventional high-scale SUSY can be easily evaded because of kinematically blocking effect.

Finally we conclude in section 5.

## 2 Implications of BICEP2 data to Inflation

In this section we report the model independent implications of BICEP2 data (together with Plank and 9-year WAMP data) to single-field inflation. These results inspired by observations on CMB temperature anisotropy and polarization provide useful information on model building of inflation, as a reliable inflation model should at least explain observable quantities as what follows.

First of all the scale of energy density during inflation is directly related to \( r \) as,

\[
V^{1/4} = \left( 24\pi^2 M_P^4 \epsilon A_s \right)^{1/4}.
\]

(2.1)

where \( A_s \) is the amplitude of the power spectrum of the curvature perturbation and \( \epsilon = r/16 \). Recall that \( A_s^{1/2} = H^2(\phi_*)/2\pi\dot{\phi} \), where \( \phi_* \) is the value of \( \phi \) when wavenumber \( k_* = 0.05 \) Mpc\(^{-1}\) crossed outside the horizon. Substituting the reported value \( A_s^{1/2} \simeq 3.089 \times 10^{-5} \) by Plank Collaboration [5] into Eq.(2.1) gives rises to [9],

\[
V^{1/4} \simeq 2 \times 10^{16} \cdot \left( \frac{r}{0.20} \right)^{1/4} \text{GeV}.
\]

(2.2)

which implies that \( V^{1/4} \sim 2 \times 10^{16} \) GeV for large value of \( r \simeq 0.1 \). If MSSM and inflation share the same origin of SUSY breaking, as we have assumed in this paper, the BICEP2 data suggests that the high-scale SUSY breaking scale in particle physics is of order \( \sqrt{F} = V^{1/4} \).
Second, in the context of slow roll inflation the spectral index $n_s$ (for scalar) and $n_t$ (for tensor) are given by,

$$n_s - 1 \simeq 2\eta - 6\epsilon, \quad n_t \simeq -2\epsilon. \tag{2.3}$$

respectively. Here $\epsilon = \frac{M_P^2}{V} (V_{,\phi}/V)^2$ and $\eta = M_P^2 V_{,\phi\phi}/V$, with subscript denoting derivative of $V$ over $\phi$. The combination of Plank and 9-year WAMP data measures the value of $n_s$ at high precision \cite{6},

$$n_s = 0.9603 \pm 0.0073 \tag{2.4}$$

which is also verified by the BICEP2 data \cite{4}. This tight bound is crucial to constrain inflation model. The reported $r$ by the Plank Collaboration with combination of 9-year WMAP data is given by $r < 0.11$ at 95% CL \cite{5,6}, whereas that reported by BICEP2 Collaboration is given by $r \simeq 0.20^{+0.06}_{-0.05}$ at 95% CL \cite{4}. This leads to tension between Plank and BICEP2 data, roughly of 3$\sigma$ level. The improved BICEP2 data later released will probably tell us whether the tension is due to uncertainty arising from foreground dust subtraction or new physics needed to reconcile this problem \cite{3}. However, if the value of $r$ is not reduced much from BICEP2 value $\sim 0.1$, it will strongly favor quadratic potential $V(\phi) \sim \frac{1}{2}m_\phi^2\phi^2 + V_0 \ [9,10,11,12]$, where $V_0$ is a constant term. The inflaton mass $m_\phi$ in the potential can be expressed in terms of slow roll parameter as,

$$m_\phi = \frac{1}{M_P} (\eta V)^{1/2} \simeq \eta^{1/2} \frac{F}{M_P} \tag{2.5}$$

where we have used Eq.(2.3) and Eq.(2.4). In Fig\ref{fig:inflaton_mass} we show uncertainty on inflaton mass defined in Eq.(2.5) due to uncertainty on $r$ and $n_s$. Typically, we have $m_\phi = \mathcal{O}(1) \times 10^{13}$ GeV for $r = 0.16 \pm 0.04$.

Third, the gravitino mass $m_{3/2}$ can be determined. The constant superpotential $W_0 = m_{3/2} M_P^2$, which is required to cancel out positive $F^2$ term in the potential so as to explain the smallness of $cc$, gives rise to

$$m_{3/2} = \frac{F}{\sqrt{3} M_P} \simeq 6m_\phi. \tag{2.6}$$

Finally, the number of e-fold that $k_*$ undergoes during inflation is given by,

$$N \simeq \int_{\phi_{end}}^{\phi_{in}} \frac{d\phi}{M_P^2 V(\phi)} \simeq \int_{x_{end}}^{x_{in}} \frac{dx}{\sqrt{2\epsilon(x)}} \tag{2.7}$$

\footnote{It is possible that the signal observed by BICEP2 attributes to the dust, see, e.g., discussions in \cite{7} and recent Plank report \cite{8} on this issue.}
Figure 1: Inflaton mass as function of $r$. The black and solid contour corresponds to central value $n_s = 0.9603$, whereas the red and dashed contours show the uncertainty on $n_s$.

where $x = \phi/M_P$. Subscript “in” and “end” corresponds to initial and end value of $x$ during inflation, respectively. For realistic inflation models, $N$ is bounded as $50 \leq N \leq 60$. If $\epsilon$ doesn’t change significantly during inflation, Eq.(2.7) can be expressed as $\Delta \phi/M_P \simeq \sqrt{2\epsilon}N \simeq \sqrt{2}N$. This is known as Lyth bound [15], which shows the need of super-Plankian excursion $\Delta \phi \sim 6M_P$ for large $r \sim 0.1$ reported by BICEP2 Collaboration.

In summary, measurements on CMB temperature anisotropy and polarization show high-scale SUSY with magnitude of order $\sim 2 \times 10^{16}$ GeV and inflaton mass $m_\phi$ of order $\sim 1 - 2 \times 10^{13}$ GeV in the course of single-field inflation.

### 3 The Simplest Hybrid Inflation

The section is devoted to the study of inflation building in the course of high-scale SUSY. We take the simplest hybrid inflation as an explicit illustration. We will show that a new linear term due to the assumption adopted in this paper (see introduction) significantly affects the choice on initial condition. Also this assumption introduces new constraints on parameters in the model, which make the simplest hybrid inflation impossible to explain the data of previous section.

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3 This bound can be evaded in small field inflation, in which $r \sim 0.1$ can be still obtained (see, e.g., [16]). The intuition is that $\epsilon$ changes significantly other than nearly is invariant during inflation.
3.1 Scalar Potential

The scalar potential in hybrid inflation is constructed from superpotential $W$, 

$$W = \kappa \Phi (\bar{\Psi} \Psi - M^2). \quad (3.1)$$

and Kahler potential $K$, 

$$K = |\Phi|^2 + |\Psi|^2 + |\bar{\Psi}|^2 + k_1 \frac{|\Phi|^4}{M_P^2}. \quad (3.2)$$

Here $\Phi$ denotes the inflaton superfield, with its lowest component inflaton field $\phi$. $\Phi$ is a singlet of standard model gauge groups $G = SU(3)_c \times SU(2)_L \times U(1)_Y$. $\Psi$ and $\bar{\Psi}$ denote waterfall superfields which are in the bi-fundamental representation of $G$.

The Kahler potential in Eq.(3.2) takes into account the non-canonical term, with $k_1$ a real coefficient. The non-canonical $k_1$ term provides the inflaton mass term as required. $M_P = 2.4 \times 10^{18}$ GeV is the reduced Plank mass, while $M$ is assumed to be far below $M_P$.

Substituting Eq.(3.1) and Eq.(3.2) into the SUGRA potential 

$$V = e^{K/M_P^2} \left[ K^{ij} D_i W \bar{D}^j W - 3 \left( \frac{|W|^2}{M_P^2} \right) \right] \quad (3.3)$$

one obtains the scalar potential of hybrid inflation. Here $K^{ij} = K^{-1}_{ij}$ is the Kahler metric and $D_i W = \partial_i W + K_i W/M_P^2$. It is well known that the history of inflation can be naturally divided into two periods. In the first one inflation usually starts from an initial value of order $M_P$ towards to $\phi_c$ which is a critical value separating the two periods. The vacuum in the first period corresponds to $\Psi = \bar{\Psi} = 0$, from which the energy density reads from Eq.(3.3) \[18] ,

$$V = \frac{1}{2} m_\phi^2 \phi^2 + \kappa^2 M_4 \left[ 1 + \gamma \frac{\phi^4}{8 M_P^4} + \frac{\kappa^2}{16 \pi^2} \ln \frac{\kappa^2 \phi^2}{2 \Lambda^2} \right] + 2 \sqrt{2} \kappa M^2 m_{3/2} \phi \cos \theta. \quad (3.4)$$

where we have defined $\Phi = \phi e^{i\theta}/\sqrt{2}$. Here $\gamma = 1 - 7k_1/2 + 2k_2^2$ and inflaton mass $m_\phi$ is given by for negative $k_1$, 

$$m_\phi = \sqrt{-k_1 \kappa M^2}/M_P. \quad (3.5)$$

Compare Eq.(3.5) with Eq.(2.5), one obtains $k_1 \simeq -\eta \simeq -0.01$. The log-term in Eq.(3.4) represents the contribution due to mass splitting in waterfall fields \[19\], with $\Lambda$ the cut-off scale. The linear term with coefficient proportional to $m_{3/2}$ arises from a constant superpotential $W_0 = m_{3/2} M_P^2$.

\[4\]Alternatively, $G$ can be extended to include a local $U(1)_{B-L}$ symmetry so as to explain leptogenesis.
added to $W$, which is needed to cancel out positive contribution to energy density due to SUSY breaking, and explain the smallness of $cc$.

When $\phi$ approaches to $\phi_c = \sqrt{2} M$, $\phi$ becomes massless as shown from its mass squared $m_\phi^2 = -4\kappa^2 M^2 + 2\kappa^2 \phi^2$. After the time when $\phi$ is below $\phi_c$, $\Psi$ starts to roll towards to its global minimum value $\Psi = M$ from $\Psi = 0$, which is known as the second period of inflation. The evaluation of field $\phi$ (including angular component $\theta$) and $\Psi$ during each period is determined by their equations of motion,

$$3H \dot{\phi} \simeq m_\phi^2 \phi + 2\sqrt{2}\kappa M^2 m_3/2 \cos \theta,$$

$$3H \dot{\theta} \simeq 2\sqrt{2}\kappa M^2 m_3/2 \frac{\sin \theta}{\phi},$$

$$3H \dot{\Psi} \simeq -2\kappa^2 (2 M^2 - \phi^2) \Psi.$$  (3.6)

where $H$ the Hubble constant is subject to the Friedmann constraint

$$H^2 = \frac{8\pi V}{3 M_P^2}.  \quad (3.7)$$

The time for each period is controlled by the magnitude of $m_\phi$ or $|m_\Psi|$ relative to Hubble constant $H$. As pointed out in [17], the second period is very short in compared with the first one for wide ranges of parameter choices. Substituting Eq.(3.7) into the last equation in Eq.(3.6), we obtain the constraint for such property,

$$10^{-4} < \kappa < O(1).  \quad (3.8)$$

In the next subsection, we will discuss in more details the initial conditions on inflaton field and the field value of $\phi$ when inflation ends.

### 3.2 Initial Conditions

The inflation usually begins at some field value $\phi_{in}$ near Plank scale. The choice on $\phi_{in}$ is subtle when the inflaton potential has either a few local minimums at $\phi_{min}$s, or local maximum at $\phi_{max}$s. If one adopts $\phi_{in}$ bigger than $\phi_{min}$, inflaton is probably trapped at these local minimums of inflaton potential along the trajectory, which leads to inflation with insufficient e-fold number $N \sim 50 - 60$. In order to avoid this, one should choose $\phi_{in} < \min\{\phi_{min}\}$. On the other hand, one wants that inflation proceeds with exactly decreasing $\phi$. This is only allowed if $\phi_{in}$ is less than $\min\{\phi_{max}\}$. In other words, we should impose the initial condition

$$\phi_{in} < \min\{\phi_{max}, \phi_{min}\}.  \quad (3.9)$$
In Fig. 2 we show how extremes in $V$ depend on $\cos \theta$ and $\kappa$ by evaluating $\sqrt{2\epsilon} = V_{,\phi}/V$. The sign of $V_{,\phi}/V$ changes when $x \sim 0.10$ for $\cos \theta = -0.002$ and $x \sim 0.15$ for $\cos \theta = -0.003$. This implies that $\cos \theta \simeq 0$ for realistic inflation. Otherwise, $\phi_{in} \ll M_P$, which is too small to provide enough e-fold number $N$. This observation has been noted in [18] for $m_{3/2}$ of order electroweak scale, and further verified for larger value of $m_{3/2} \sim 10^{13}$ GeV. With initial value $\theta \simeq \pi/2$, the initial value $\phi_{in}$ can be chosen in wide range, as shown in Fig. 2. The evaluation of $\phi$ from $\phi_{in}$ is the same as original hybrid model because of absence of linear term in the first equation in Eq. (3.6).

In this sense, inflation mainly ends at field value $\phi_{end} = \sqrt{2} M$.

The e-fold number $N$ produced during inflation and $n_s$ can be both estimated in terms of slow roll parameters $\epsilon$ and $\eta$ in the model, which are given by respectively,

\[
\eta(x) \equiv M_P^2 \frac{V_{,\phi}}{V} \simeq -k_1 + \frac{3}{2} x^2 - \frac{\kappa^2}{8\pi^2} \frac{1}{x^2},
\]

\[
\epsilon(x) \equiv \frac{M_P^2}{2} \left( \frac{V_{,\phi}}{V} \right)^2 \simeq \frac{1}{2} \left( \sqrt{\frac{8}{3}} \cos \theta - k_1 x + \frac{1}{2} x^3 + \frac{\kappa^2}{8\pi^2} \frac{1}{x} \right)^2.
\]  

With $k_1 \simeq -0.01$ and $\cos \theta \simeq 0$, $\eta(x)$ and $\epsilon(x)$ mainly depend on parameter $\kappa$.

Fig. 3 shows the bound on $\kappa$ for two choices of $x_{in} = 0.3$ (red curve) and 0.5 (black curve) respectively. Note that $x$ is constrained from slow roll condition $| \eta | < 1$. Given the range shown in Eq. (3.8) for $\kappa$, $x$ should be below unity, which implies that large field inflation is excluded under our assumption. Fig. 3 shows that $\kappa \sim 0.1$ for $N \sim 50 - 60$.

In Fig. 4 we show how $n_s$ changes for two typical choices of $\kappa$ subtracted from Fig. 3. It clearly indicates that for observed value of $n_s$, $r$ is of order $\sim 10^{-5}$. The simplest hybrid inflation can
Figure 3: \( N \) as function of \( \kappa \) for initial value \( x_{in} = 0.3 \) (red), 0.5 (black) respectively. Dotted lines represent the uncertainty on experimental value. Slow roll condition \( |\eta| < 1 \) leads to the bound \( \phi_{in} \leq M_P \).

Figure 4: \( n_s \) as function of \( \kappa \) for initial value \( x_{in} = 0.3 \) (red), 0.5 (black) respectively. Note that \( x_{end} \geq \phi_c/M_P \). Dotted lines represent the uncertainty on experimental value.
provide large e-folds number $N \sim 50 - 60$ and small $r << 0.1$, or vice versa. Nevertheless, large $N$ and large $r$ as required can not be induced at the same time. So the simplest hybrid inflation doesn’t explain BICEP2 data in the course of high-scale SUSY. In the next section, we focus on reheating after (or during) inflation.

4 Reheating in High-Scale SUSY

When inflation approaches to end, the conversion of energy in inflaton sector to MSSM matters begins. The efficiency of energy transfer depends on how inflaton is coupled to MSSM matters, the magnitude of their couplings, and the superpartner spectrum. In general, the ways of energy transfer include perturbative decay and non-perturbative decay of inflaton. The later way is also known as preheating [13, 21, 22]. The conditions between these two ways of energy transfer are rather different. In the later case, parameter resonance requires a quartic interaction term $\sim g^2 \phi^2 \chi^2$ with large magnitude of $g$. This only happens if one allows renormalizable superpotential term of mass dimension 4 [23],

$$
R_\phi = +1 : \Phi H_u H_d,
$$

$$
R_\phi = -1 : \Phi H_u L.
$$

(4.1)

where $\Phi$ is the inflaton superfield, and $H_{u,d}$ are Higgs doublet superfields. $R_\phi$ denotes the $R$-parity of inflaton, which is useful to keep dark matter stable. In contrast, in SM the inflaton couples to SM chiral fermions and gauge bosons in terms of non-renormalizable interactions of mass dimension 5. Quartic term above doesn’t exist in SM, and therefore the way of energy transfer in SM is perturbative decay. In what follows, we consider these two ways separately.

4.1 Perturbative Decay

As briefly mentioned above, perturbative decay happens either when there is no renormalizable interaction in Eq.(4.1) or the quartic coupling constant is tiny. Instead, the inflaton only decays to SM matters via five-dimensional operators such as

$$
\{ \frac{\phi}{M} F_{\mu\nu} F^{\mu\nu}, \frac{\phi}{M} \phi (H \bar{q}_L) q_R, \cdots \}.
$$

(4.2)

Here $F_{\mu\nu}$s refer to strengths of SM gauge fields, $q$s refer to SM fermions and $M$ represents the mass scale appearing in the five-dimensional operators. The plasma will be MSSM-like if the reheating

\footnote{For recent attempts to fix this problem, see Refs.\[11\,20\].}

\footnote{For reviews, see, e.g., \[13\,14\].}
temperature is larger than the typical scale of superpartner mass, $m_0$. Otherwise, the plasma is actually SM-like.

Now we calculate the reheating temperature. We organize the decay width of inflaton to SM particles as,

$$\Gamma_d \equiv \frac{\lambda}{16\pi^2} \left( \frac{m_{\phi}}{M} \right)^2 m_{\phi}$$  \hspace{1cm} (4.3)

We simply take $M = M_P$ but leave $\lambda$ as a free parameter. The thermal equilibrium of relativistic plasma is dominated by $\Gamma$, the rate for SM inelastic scatterings of $2 \rightarrow 3$ processes [24]. $\Gamma$ is related to $\Gamma_d$ as,

$$\Gamma \sim \alpha^3 \left( \frac{M_P}{m_{\phi}} \right) \Gamma_d,$$  \hspace{1cm} (4.4)

where $\alpha \sim 1/30$ is the SM fine structure constant. Therefore the reheating temperature $T_R$ for the perturbative decay isn’t equal to the conventional one, $T_{rth}$ defined as $T_{rth} \simeq 0.3 \times \left( \frac{100}{g_*} \right)^{1/4} \sqrt{\frac{1}{4dM_P}}$, $g_*$ denotes the number of relativistic number. Instead, in terms of Eq. (4.4) we have

$$T_R \simeq \alpha^{3/2} \cdot \left( \frac{2\pi N_c M_P}{0.09 m_{\phi}} \right)^{1/2} \cdot \left( \frac{g_{\text{susy}}}{g_{\text{sm}}^2} \right)^{1/4} \cdot T_{rth} \simeq 0.01 \cdot \left( \frac{g_{\text{susy}}}{g_{\text{sm}}^2} \right)^{1/4} \cdot \left( \frac{\lambda^{1/2}}{4\pi} \right) \cdot m_{\phi}$$  \hspace{1cm} (4.5)

where $N_c$ denote the quantum numbers of SM gauge groups. Note that Eq. (4.5) is valid for $m_0 < T_R$, which implies that the reheating temperature can serve as the upper bound on $m_0$.

In Fig 5, we show reheating temperature $T_R$ as function of inflaton mass. Note that $\lambda$ captures the magnitude of coupling between inflaton and “mediate ” field, which also couples to the SM matter and gauge fields. Here a few comments are in order. (1), For $\lambda < 10^{-3}$, $T_R$ is below the lower bound $\sim 1 \times 10^9$ GeV required by thermal leptogenesis in the whole range of $m_{\phi}$. (2), Since $T_R$ is the upper bound on superpartner mass spectrum $m_0$, one finds that $m_0$ is upper bounded as $1 \text{ TeV} << m_0 < 10^7 \text{ TeV}$ for the case of perturbative decay. (3), As $T_R$ is far below the gravitino mass of Eq. (2.6), there is no overproduction problem of gravitino in high-scale SUSY. Superheavy gravitino mass of order $\sim 10^{14} \text{ GeV}$ kinetically blocks its production in the thermal bath.

### 4.2 Non-perturbative Decay

If it admits renormalizable superpotential Eq. (4.1), there exists quartic interaction between inflaton and its decay products $\chi_i$, Non-perturbative decay can happen in wide range of parameter space
for potential of type\footnote{Assuming inflaton and MSSM matters share the same origin of SUSY breaking, inflaton mass is dynamical induced by SUSY breaking. In this sense, the mass term $m_\phi^2 \phi^2$ is a soft SUSY-breaking term other than arises from SUSY tree-level mass superpotential $\sim m\Phi\Phi$. Consequently, there is no cubic interaction $\phi\chi^2$ in compared with earlier discussions in \cite{25,27}.},

$$V(\phi, \chi) = \frac{1}{2}m_\phi^2 \phi^2 + g^2 \phi^2 \chi_i^2 + m_{\chi_i}^2 \chi_i^2. \quad (4.6)$$

where the inflaton mass term is included and $m_{\chi_i}$ is the mass for $\chi_i$. We would like to mention that $m_{\chi_i}$ include soft SUSY breaking contribution of order $\sim m_0$ and dynamical mass $\sim \lambda^{1/2} \langle \phi \rangle$ induced by VEV of flat direction $\phi$ \cite{25} through the quartic interaction \cite{27},

$$V(\chi, \phi) = \lambda \chi_i^2 \phi^2. \quad (4.7)$$

where $\lambda$ is the quartic coupling constant. The magnitude of $\langle \phi \rangle$ is determined by the self-interaction potential for flat direction $V(\phi)$.

Now we consider the potential for flat direction. $V(\phi)$ includes soft breaking mass, Hubble parameter induced term and high dimensional operators,

$$V(\phi) \simeq (m_0^2 + c_H H^2) \phi^2 + c_6 \frac{\phi^6}{M^4} + \cdots \quad (4.8)$$

where $c_H$ is real coefficient. Since $m_0$ isn’t far beneath the Hubble constant $H \sim m_\phi$ at the beginning of inflation, VEV $\langle \phi \rangle$ depends on the sign of $c_H$, which can be either positive or negative \cite{25,26}. In particular, $\langle \phi \rangle = 0$ for the case of either positive $c_H$ or negative $c_H$ but with $|c_H| << 1$.\footnote{Assuming inflaton and MSSM matters share the same origin of SUSY breaking, inflaton mass is dynamical induced by SUSY breaking. In this sense, the mass term $m_0^2 \phi^2$ is a soft SUSY-breaking term other than arises from SUSY tree-level mass superpotential $\sim m\Phi\Phi$. Consequently, there is no cubic interaction $\phi\chi^2$ in compared with earlier discussions in \cite{25,27}.}

Figure 5: Reheating temperature (also the upper bound on $m_0$) as function of inflaton mass.
It implies that SM gauge symmetry is unbroken during the whole history of early universe. On the other hand, \( \langle \varphi \rangle \neq 0 \) for negative \( c_H \) but with \( |c_H| > 1 \). It implies that SM gauge symmetry is broken in the early universe, with gauge boson mass of order \( \langle \varphi \rangle \), then restored after the epoch of reheating.

The potentials we define in Eq.(4.6) to Eq.(4.7) are rather general, which can be applied to both cases in Eq.(4.1). To discuss the condition for parameter resonance, one starts with the modified Klein-Gordon equation for Fourier modes \( \chi \). Whether WKB approximation is viable for the study can be analyzed in term of a quantity \( R \) defined as [13],

\[
R \equiv \frac{\dot{\omega}_k}{\omega_k^2} \quad \omega^2 = k^2/a^2 + g^2 \langle \phi(t) \rangle^2.
\]  

(4.9)

where \( \dot{\omega}_k \) refers to derivative over time and \( \omega \) is the frequency, with \( a \) the expansion factor and \( k \) the momentum. If \( |R| \ll 1 \), the WKB approximation is valid, the produced particle number of \( \chi \) doesn’t grow in this case. If \( |R| > 1 \) instead, the WKB approximation isn’t valid, which leads to significant production of \( \chi \). In long wavelengths limit, this constraint is given by \( 8 \),

\[
m^2_{\chi} \simeq m_0^2 + \lambda \langle \varphi \rangle^2 < g^2 \langle \phi(t) \rangle^2.
\]  

(4.10)

where we have used Eq.(4.6) and Eq.(4.7). Moreover, in order to keep that the parameter resonance isn’t spoiled by expansion, an additional constraint must be imposed,

\[
q \equiv g^2 \bar{\phi}^2(t)/4m_\phi^2 >> 1
\]  

(4.11)

where \( \bar{\phi}(t) \sim M_P \) refers to the amplitude of inflaton oscillations. For more details, we refer to reader to [13] and references therein.

Fig 6 shows reheating temperature for the case in which \( \langle \varphi \rangle \neq 0 \). In this case, SM gauge bosons are massive and the rate for thermal equilibrium \( \Gamma \) [27] depends on its magnitude relative to \( m_0 \). For \( m_0 < \Gamma_d \), \( \Gamma_R \) depends on both \( g \) and \( \langle \varphi \rangle \), whereas it mainly depends on \( \langle \varphi \rangle \) for \( m_0 > \Gamma_d \). In this figure, we take \( m_\phi = 10^{13} \) GeV and \( \bar{\phi} = M_P \). The bounds on \( q \) are due to a few considerations. The first one is that condition Eq.(4.11) from parameter resonance requires \( g >> 10^{-5} \). The second one is that overproduction problem [28] of gravitino in high-scale SUSY can be kinematically blocked if \( g < 10^{-2} \) such that non-perturbatively induced mass during parameter resonance is beneath \( m_{3/2} \). The bound on \( \lambda \langle \varphi \rangle \) arises from condition Eq.(4.10) which shows \( \lambda \langle \varphi \rangle^2 < g^2 M_P^2 \).

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8There is a coefficient of order one in front of \( m_\phi \) for either \( R_\phi = 1 \) or \( R_\phi = -1 \). Here we simply take it equal to unity.

9 Ref. [29] provides an example how gravitino problem in high-scale SUSY is evaded in the context of mini-Split SUSY. In comparison with [29], gravitino mass is far heavier in this paper, and kinematically blocking is the solution to the overproduction of gravitino.
Figure 6: Reheating temperature as function of $m_0$ for $m_\phi = 10^{13}$ GeV and $\langle \varphi \rangle \neq 0$. (Left: $m_0 < \Gamma_d$ with $\langle \varphi \rangle = \mathcal{M}_P$, Right: $m_0 > \Gamma_d$ with $\lambda = 10^{-3}$ and $\langle \varphi \rangle = 0.01 \mathcal{M}_P$). In the Right panel, $\langle \varphi \rangle \geq 0.1 \mathcal{M}_P$ is excluded by condition Eq.(4.10) for $\lambda \geq 10^{-3}$. The bounds on $g$ and $m_0$ are explained in the text.

The Left panel in Fig.6 shows that in the range $10^5$ GeV < $m_0 < 1.5 \times 10^7$ GeV reheating temperature $T_R \geq 10^9$ GeV in the allowed range of $g$ for $\langle \varphi \rangle = \mathcal{M}_P$. With modifying $\langle \varphi \rangle < \mathcal{M}_P$, 

$$T_R \to \left( \frac{\langle \varphi \rangle}{\mathcal{M}_P} \right)^{-1} T_R,$$

which is always above the value required by thermal leptogenesis.

The Right panel in Fig.6 shows that in the range $10^{11}$ GeV < $m_0 < 4 \times 10^{13}$ GeV reheating temperature $10^{11}$ GeV \leq T_R \leq 10^{13}$ GeV if $\langle \varphi \rangle = 10^{-2} \mathcal{M}_P$, and changes similarly to Eq.(4.12) for modifying $\langle \varphi \rangle$. This implies that $T_R$ is also always above the the value required by thermal leptogenesis. Note that $\langle \varphi \rangle \geq 10^{-1} \mathcal{M}_P$ is excluded by condition Eq.(4.10) from parameter resonance for $\lambda \simeq 10^{-3}$.

Fig.7 shows reheating temperature for the case in which $\langle \varphi \rangle = 0$. In this case, SM gauge symmetries are unbroken in the epoch of reheating. Thermalization cannot occur before the inflaton decay has completed. Due to $\langle \varphi \rangle = 0$ the $2 \to 3$ scatterings are already efficient when $H \simeq \Gamma_d$. The reheat temperature in this case is given by the standard expression: $T_R \simeq 0.3 \sqrt{T_d \mathcal{M}_P}$. However, $\Gamma_d$ calculated via renormalizable couplings is rather different from Eq.(4.3) calculated via non-renormazible couplings. This difference between Fig.5 and Fig.7 is obvious. Typically, we have $T_R \geq 10^{10}$ GeV in non-perturbative decay into MSSM and $T_R \leq 10^{10}$ GeV in perturbative decay into SM.

Both Fig.7 and Fig.6 show that $T_R$ is above $\sim 10^9$ GeV but beneath $m_{3/2}$ in a wide range of parameter space. Due to kinematically blocking effect this evades the overproduction of gravitino.
Figure 7: Reheating temperature as function of $g$ for $\langle \varphi \rangle = 0$, with the definition $\Gamma_d \equiv g^2 m_{\phi}/8\pi$.

in conventional high-scale SUSY.

5 Conclusions

High-scale SUSY is more attractive in compared with low-scale SUSY in light of both the LHC and BICEP2 data. In this paper, we discuss implications of high-scale SUSY to the early universe. In particular, we assume that the inflation and MSSM share the same origin of SUSY breaking, and derive model independent consequences based on this assumption. We find that the reheating temperature for superpartner mass spectrum $m_0 > \mathcal{O}(100)$ TeV might be beneath the value required by thermal leptogenesis if inflaton decays to its products perturbatively but above it if non-perturbatively instead. We also observe that problem of gravitino overproduction can be evaded through kinematically blocking in a wide range of parameter space in the later way.

As an illustration for model building of inflation in the course of high-scale SUSY, in section 3 we revise the simplest hybrid inflation that include a new linear term for inflaton with coefficient proportional to $m_{3/2}$. It is shown that this term significantly affects the choices on initial condition of inflaton fields. We find that with the assumption we follow the simplest hybrid inflation is impossible to explain all experimental data of section 2 simultaneously.

There are a few interesting issues along this line. For example, dark matter and the collider physics in such high-scale SUSY deserves detailed study \cite{30}. The early universe with high-scale SUSY induced by $r$ which is far beneath the BICEP2 value also deserves detailed study.
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