

Reheat temperature in supersymmetric hybrid inflation models

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The allowed range of parameters for supersymmetric hybrid inflation and its extensions are investigated. The lower bound on the reheat temperature T_r in these models with hierarchical right-handed neutrinos is found to be 3×10^7 GeV. (T_r as low as 100 GeV is possible for quasidegenerate right-handed neutrinos.) We also present revised estimates for the scalar spectral index and the symmetry breaking scale associated with inflation.

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I. INTRODUCTION

In supersymmetric (SUSY) hybrid inflation models, inflation is associated with the breaking of a gauge symmetry G to the minimal supersymmetric standard model (MSSM). The symmetry breaking scale M is fixed by the amplitude of the primordial curvature perturbation, and turns out to be of order 10^{16} GeV, remarkably close, if not identical, to the supersymmetric grand unification scale [1,2].

A particularly simple and compelling example of G is provided by the standard model gauge group supplemented by a gauged $U(1)_{B-L}$ symmetry which requires, from anomaly cancellations, the presence of three right-handed neutrinos. Other, more elaborate examples of G in which SUSY hybrid inflation [1] and its extensions, namely, shifted hybrid inflation [3] and smooth hybrid inflation [4], have been implemented include $SO(10)$ [5] and its various subgroups. In the hybrid and shifted hybrid inflation models, the scalar spectral index $n_s \geq 0.98$, whereas in the smooth case $n_s \geq 0.97$. Furthermore, $dn_s/d \ln k \leq 10^{-3}$, while the tensor to scalar ratio r turns out to be of order 10^{-4} or less [6].

The main focus of this paper is to estimate the reheat temperature that is required to generate sufficient lepton asymmetry following hybrid, shifted or smooth hybrid inflation. Although reheating and leptogenesis in these models has previously been addressed [7–10], the soft SUSY breaking terms and their impact on inflation were not adequately included. It turns out that the new terms are significant for estimating the minimum reheat temperature in these models. We also present revised estimates for the symmetry breaking scale associated with inflation as well as the scalar spectral index.

The layout of the paper is as follows. In Sec. II we provide a brief review of SUSY hybrid inflation and its extensions. We compute the allowed range of the dimensionless coupling in the superpotential and the dependence of the spectral index on this coupling, in the presence of canonical SUGRA corrections and soft SUSY violating

terms. In Sec. III we investigate reheating and the generation of matter following inflation in these models. For hierarchical right-handed neutrinos, we obtain a lower bound on the reheat temperature from the observed baryon asymmetry. We also briefly comment on the resolution of the MSSM μ problem and its impact on the reheat temperature.

II. SUSY HYBRID INFLATION MODELS

The SUSY hybrid inflation model [1] is realized by the renormalizable superpotential [11]

$$W_1 = \kappa S(\Phi\bar{\Phi} - M^2) \quad (1)$$

where $\Phi(\bar{\Phi})$ denote a conjugate pair of superfields transforming as nontrivial representations of some gauge group G , S is a gauge singlet superfield, and $\kappa (>0)$ is a dimensionless coupling. A suitable $U(1)$ R-symmetry, under which W_1 and S transform the same way, ensures the uniqueness of this superpotential at the renormalizable level [1]. In the absence of supersymmetry breaking, the potential energy minimum corresponds to nonzero (and equal in magnitude) vacuum expectation values (vevs) ($= M$) for the scalar components in Φ and $\bar{\Phi}$, while the vev of S is zero. (We use the same notation for superfields and their scalar components.) Thus, G is broken to some subgroup H which, in many interesting models, coincides with the MSSM gauge group.

In order to realize inflation, the scalar fields Φ , $\bar{\Phi}$, S must be displaced from their present minima. For $|S| > M$, the Φ , $\bar{\Phi}$ vevs both vanish so that the gauge symmetry is restored, and the tree level potential energy density $\kappa^2 M^4$ dominates the universe. With supersymmetry thus broken, there are radiative corrections from the $\Phi - \bar{\Phi}$ supermultiplets that provide logarithmic corrections to the potential, with additional contributions to the inflationary potential arising from $N = 1$ supergravity.

With a minimal Kahler potential, one contribution to the inflationary potential is given by [11–13]

$$V_{\text{SUGRA}} = \kappa^2 M^4 \left[\frac{|S|^4}{2m_p^4} + \dots \right], \quad (2)$$

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where $m_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. There are additional contributions to the potential arising from the soft SUSY breaking terms. In $N = 1$ SUGRA these include the universal scalar masses equal to $m_{3/2}$ (\sim TeV), the gravitino mass. However, their effect on the inflationary scenario is negligible, as discussed below. The more important soft term, ignored as far as we can tell in all earlier calculations, is $(2 - A)m_{3/2}\kappa M^2 S(+\text{H.c.})$. For convenience, we write this as $am_{3/2}\kappa M^2|S|$, where $a \equiv 2|2 - A| \cos[\arg S + \arg(2 - A)]$. The effective potential is approximately given by the radiative corrections [1] plus the leading SUGRA correction $\kappa^2 M^4 |S|^4 / 2m_P^4$ and this soft term:

$$V_1 = \kappa^2 M^4 \left\{ 1 + \frac{\kappa^2 \mathcal{N}}{32\pi^2} \left[2 \ln \frac{\kappa^2 |S|^2}{\Lambda^2} + (z+1)^2 \ln(1+z^{-1}) + (z-1)^2 \ln(1-z^{-1}) \right] + \frac{|S|^4}{2m_P^4} \right\} + am_{3/2}\kappa M^2 |S|, \quad (3)$$

where $z \equiv |S|^2/M^2$, \mathcal{N} is the dimensionality of the Φ , $\bar{\Phi}$ representations, and Λ is a renormalization mass scale. We perform our numerical calculations using this potential, taking $|am_{3/2}| = 1$ TeV, and using the well-known equations in the slow roll approximation: The number of e-folds after the comoving scale l has crossed the horizon is given by

$$N_l = \frac{1}{m_P^2} \int_{\sigma_f}^{\sigma_l} \frac{V d\sigma}{V'}. \quad (4)$$

Here $\sigma \equiv \sqrt{2}|S|$ is the normalized real scalar field, σ_l is the value of the field at the comoving scale l , and σ_f is the value of the field at the end of inflation. The amplitude of the curvature perturbation \mathcal{R} is given by

$$\mathcal{R} = \frac{1}{2\sqrt{3}\pi m_P^3} \frac{V^{3/2}}{|V'|}, \quad (5)$$

which we evaluate for the comoving wave number $k_0 \equiv 0.002 \text{ Mpc}^{-1}$.

It is instructive to discuss small and large κ limits of Eq. (3). For $\kappa \gg 10^{-3}$, Eq. (3) becomes

$$V_1 \simeq \kappa^2 M^4 \left[1 + \frac{\kappa^2 \mathcal{N}}{32\pi^2} 2 \ln \frac{\kappa^2 |S|^2}{\Lambda^2} + \frac{|S|^4}{2m_P^4} \right], \quad (6)$$

to a good approximation. This potential has been analyzed in [6,13], and the presence of the SUGRA correction was shown to lead to a blue spectrum for $\kappa \gtrsim 0.06/\sqrt{\mathcal{N}}$.

For $\kappa \ll 10^{-3}$, $|S_0| \simeq M$ where S_0 is the value of the field at k_0 , i.e., $z \simeq 1$. (Note that due to the flatness of the potential the last 55 or so e-folds occur with $|S|$ close to M .) From Eqs. (3) and (5), as $z \rightarrow 1$

$$\mathcal{R} = \frac{2\sqrt{2}\pi}{\sqrt{3}m_P^3} \frac{\kappa^2 M^4}{\mathcal{N} \ln(2)\kappa^3 M + 8\pi^2 \kappa M^5/m_P^4 + 4\pi^2 am_{3/2}}. \quad (7)$$

The denominator of Eq. (7) contains the radiative, SUGRA and the soft terms, respectively. Comparing them, we see that the radiative term can be ignored for $\kappa \lesssim 10^{-4}$.

For a positive soft term ($a > 0$), the maximum value of \mathcal{R} as a function of M is found to be

$$\mathcal{R}_{\max} = \frac{1}{2^{7/10} 35^{3/2} \pi} \left(\frac{\kappa^6 m_P}{am_{3/2}} \right)^{1/5}. \quad (8)$$

Setting $\mathcal{R} \simeq 4.7 \times 10^{-5}$ [14], we find a lower bound on κ ($\simeq 10^{-5}$). For larger values of κ , there are two separate solutions of M for a given κ .

For $a < 0$, there are again two solutions, but for the solution with a lower value of M , the slope changes sign as the inflaton rolls for $\kappa \lesssim 10^{-4}$ and the inflaton gets trapped in a false vacuum. The second solution in principle allows $\kappa < 10^{-5}$, but this is not very natural since it requires a delicate cancellation between two large terms in the denominator of Eq. (7).

There is also a soft mass term $m_{3/2}^2 |S|^2$ in the potential, corresponding to an additional term $8\pi^2 m_{3/2}^2 / \kappa M$ in the denominator of Eq. (7). We have omitted this term, since it is insignificant for $\kappa \gtrsim 10^{-5}$.

The dependence of M on κ as well as the allowed range of κ is shown in Fig. 1. Note that the soft term depends on $\arg S$, so it should be checked whether $\arg S$ changes significantly during inflation. Numerically, we find that it does not, except for a range of κ around 10^{-4} . For this range, corresponding to the gray segments in the figure, if the initial value of the S field is greater than M by at least a

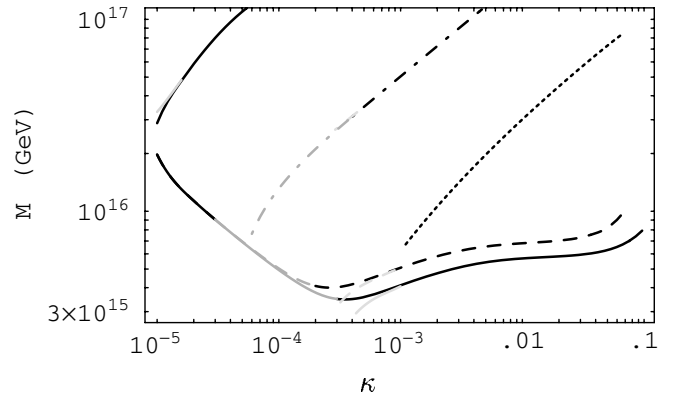


FIG. 1. The value of the symmetry breaking scale M vs the allowed range of κ , for SUSY hybrid inflation with $\mathcal{N} = 1$ (solid line), with $\mathcal{N} = 2$ (dashed line), and for shifted hybrid inflation (dot-dashed line for $M_S = m_P$, dotted for $M_S = 5 \times 10^{17}$ GeV). Light gray portions of the curves are for $a < 0$, where only the segments that do not overlap with the solutions for $a > 0$ are shown. The gray segments denote the range of κ for which the change in $\arg S$ is significant.

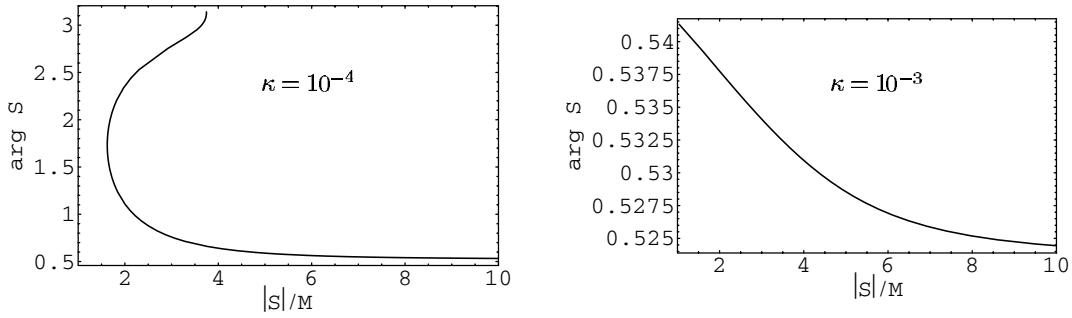


FIG. 2. Two examples of how $\arg S$ changes as S rolls down, for SUSY hybrid inflation with $\mathcal{N} = 2$. Left: $\arg S$ exceeds $\pi/2$ before the field reaches the waterfall point, and the field relaxes in a false vacuum. Right: The field reaches the waterfall point without a significant change in $\arg S$. The initial value of $\arg S = \pi/6$, $\arg(2 - A)$ is taken to be zero.

factor of 2 or so, the soft term and the slope become negative even if they were initially positive, before inflation can suitably end (Fig. 2). As mentioned above, $|S| \simeq M$ during the last 55 or so e-folds, so strictly speaking this range is not excluded, although the required initial conditions may look contrived.

Qualitatively, the duration of inflation is given by N/H , where N denotes the total number of e-folds, and the Hubble parameter $H \propto \kappa M^2/m_P$. The rate of change in $\arg S$ is $\propto (m_{3/2} m_P)/|S|$. Therefore one expects that $\arg S$ would change significantly if $\kappa M^3/N \lesssim m_{3/2} m_P^2$. For the range of κ where the radiative term dominates, $N \propto \kappa^{-2}$ [13]. However, for $\kappa \lesssim 10^{-4}$, the soft term dominates the slope of the potential, and Eq. (7) then has solutions with higher values of the slope, and the duration of inflation is shorter. Consequently, $\arg S$ stays fixed also in the left segments of the curves.

The dependence of n_s on κ is displayed in Fig. 3. The segment with $n_s > 1$ for small κ corresponds to the solution with a high symmetry breaking scale. The running of the spectral index is negligible, with $dn_s/d \ln k \lesssim 10^{-3}$.

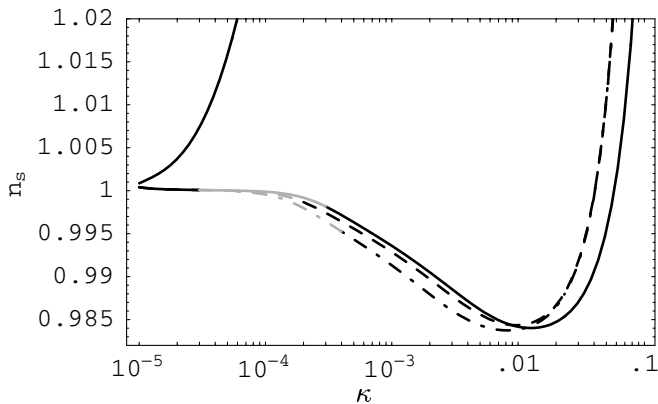


FIG. 3. The spectral index n_s vs the allowed range of κ , for SUSY hybrid inflation with $\mathcal{N} = 1$ (solid line), with $\mathcal{N} = 2$ (dashed line), and for shifted hybrid inflation with $M_S = m_P$ (dot-dashed line). The gray segments denote the range of κ for which the change in $\arg S$ is significant.

The experimental data disfavor n_s values in excess of unity on smaller scales (say $k \lesssim 0.05 \text{ Mpc}^{-1}$), which leads us to restrict ourselves to $\kappa \lesssim 0.1/\sqrt{\mathcal{N}}$ for $n_s \leq 1.04$.¹ Thus, the vacuum energy density during inflation is considerably smaller than the symmetry breaking scale. Indeed, the tensor to scalar ratio $r \lesssim 10^{-4}$.

The inflationary scenario based on the superpotential W_1 in Eq. (1) has the characteristic feature that the end of inflation essentially coincides with the gauge symmetry breaking. Thus, modifications should be made to W_1 if the breaking of G to H leads to the appearance of topological defects such as monopoles, strings or domain walls. As shown in [3], one simple resolution of the topological defects problem is achieved by supplementing W_1 with a nonrenormalizable term:

$$W_2 = \kappa S(\overline{\Phi}\Phi - v^2) - \frac{S(\overline{\Phi}\Phi)^2}{M_S^2}, \quad (9)$$

where v is comparable to the SUSY grand unified theory (GUT) scale $M_{\text{GUT}} \simeq 2 \times 10^{16} \text{ GeV}$ and M_S is an effective cutoff scale. The dimensionless coefficient of the nonrenormalizable term is absorbed in M_S . The presence of the nonrenormalizable term enables an inflationary trajectory along which the gauge symmetry is broken. Thus, in this “shifted” hybrid inflation model the topological defects are inflated away.

The inflationary potential is similar to Eq. (3) [3]:

$$V_2 = \kappa^2 m^4 \left\{ 1 + \frac{\kappa^2}{16\pi^2} \left[2 \ln \frac{\kappa^2 |S|^2}{\Lambda^2} + (z+1)^2 \ln(1+z^{-1}) + (z-1)^2 \ln(1-z^{-1}) \right] + \frac{|S|^4}{2m_P^4} \right\} + am_{3/2} \kappa v^2 |S|. \quad (10)$$

Here $m^2 = v^2(1/4\xi - 1)$ with $\xi = v^2/\kappa M_S^2$, $z \equiv$

¹Larger values of κ may be allowed in models where dissipative effects are significant. Such effects become important for large values of κ , provided the inflaton also has strong couplings to matter fields [15].

$2|S|^2/m^2$, and $2 - A$ is replaced by $2 - A + A/2\xi$ in the expression for a . The slow roll parameters (and therefore n_s , $dn_s/d\ln k$, and r) are similar to the SUSY hybrid inflation model (Fig. 3).

The vev M at the SUSY minimum is given by [3]

$$\left(\frac{M}{v}\right)^2 = \frac{1}{2\xi}(1 - \sqrt{1 - 4\xi}), \quad (11)$$

and is $\sim 10^{16}$ – 10^{17} GeV depending on κ and M_S . The system follows the inflationary trajectory for $1/7.2 < \xi < 1/4$ [3], which is satisfied for $\kappa \gtrsim 10^{-5}$ if the effective cutoff scale $M_S = m_p$. For lower values of M_S , the inflationary trajectory is followed only for higher values of κ , and M is lower for a given κ (Fig. 1).

A variation on these inflationary scenarios is obtained by imposing a Z_2 symmetry on the superpotential, so that only even powers of the combination $\Phi\bar{\Phi}$ are allowed [4,6,16]:

$$W_3 = S\left(-v^2 + \frac{(\Phi\bar{\Phi})^2}{M_S^2}\right), \quad (12)$$

where the dimensionless parameter κ is absorbed in v . The resulting scalar potential possesses two (symmetric) valleys of local minima which are suitable for inflation and along which the GUT symmetry is broken. As in the case of shifted hybrid inflation, potential problems associated with topological defects are avoided.

The vev M at the SUSY minimum is given by $(vM_S)^{1/2}$. For $|S| \gg M$, the inflationary potential is

$$V_3 \approx v^4 \left[1 - \frac{1}{54} \frac{M^4}{|S|^4} + \frac{|S|^4}{2m_p^4} \right], \quad (13)$$

where the last term arises from the canonical SUGRA correction. The soft terms in this case do not have a significant effect on the inflationary dynamics. In the absence of the SUGRA correction $n_s \simeq 0.97$ [4]. The SUGRA correction raises n_s to above unity for $M \gtrsim 1.5 \times$

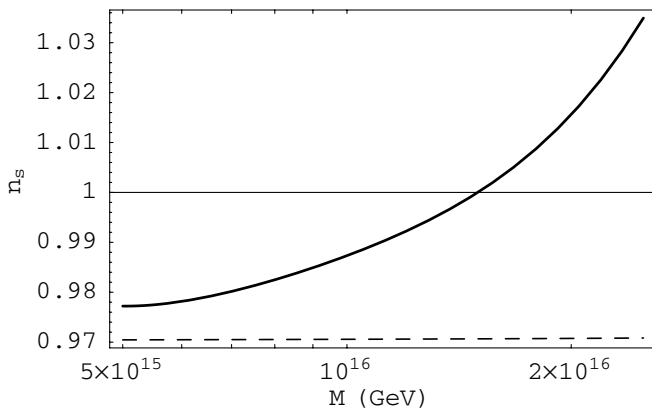


FIG. 4. The spectral index n_s as a function of the gauge symmetry breaking scale M for smooth hybrid inflation (dashed line—without SUGRA correction, solid line—with SUGRA correction).

10^{16} GeV, as shown in Fig. 4 (this figure is slightly different from the figure published in [6], due to a computational error in the latter).

III. REHEAT TEMPERATURE AND THE GRAVITINO CONSTRAINT

An important constraint on SUSY hybrid inflation models arises from considering the reheat temperature T_r after inflation, taking into account the gravitino problem which requires that $T_r \lesssim 10^6$ – 10^{11} GeV [17]. This constraint on T_r depends on the SUSY breaking mechanism and the gravitino mass $m_{3/2}$. For gravity mediated SUSY breaking models with unstable gravitinos of mass $m_{3/2} \simeq 0.1$ – 1 TeV, $T_r \lesssim 10^6$ – 10^9 GeV [18], while $T_r \lesssim 10^{10}$ GeV for stable gravitinos [19]. In gauge mediated models the reheat temperature is generally more severely constrained, although $T_r \sim 10^9$ – 10^{10} GeV is possible for $m_{3/2} \simeq 5$ – 100 GeV [20]. Finally, the anomaly mediated symmetry breaking scenario may allow gravitino masses much heavier than a TeV, thus accommodating a reheat temperature as high as 10^{11} GeV [21].

After the end of inflation in the models discussed in Sec. II, the fields fall toward the SUSY vacuum and perform damped oscillations about it. The vevs of $\bar{\Phi}$, Φ along their right-handed neutrino components $\bar{\nu}_H^c$, ν_H^c break the gauge symmetry. The oscillating system, which we collectively denote as χ , consists of the two complex scalar fields $(\delta\bar{\nu}_H^c + \delta\nu_H^c)/\sqrt{2}$ (where $\delta\bar{\nu}_H^c$, $\delta\nu_H^c$ are the deviations of $\bar{\nu}_H^c$, ν_H^c from M) and S , with equal mass m_χ .

We assume here that the inflaton χ decays predominantly into right-handed neutrino superfields N_i , via the superpotential coupling $(1/m_p)\gamma_{ij}\bar{\phi}\bar{\phi}N_iN_j$ or $\gamma_{ij}\bar{\phi}\bar{\phi}N_iN_j$, where i, j are family indices (see later for a different scenario connected to the resolution of the MSSM μ problem). Their subsequent out of equilibrium decay to lepton and Higgs superfields generates lepton asymmetry, which is then partially converted into the observed baryon asymmetry by sphaleron effects [22]. The right-handed neutrinos, as shown below, can be heavy compared to the reheat temperature T_r . Without this assumption, the constraints to generate sufficient lepton asymmetry would be more stringent [23].

GUTs typically relate the Dirac neutrino masses to that of the quarks or charged leptons. It is therefore reasonable to assume the Dirac masses are hierarchical. The low-energy neutrino data indicates that the right-handed neutrinos in this case will also be hierarchical in general. As discussed in Ref. [24], setting the Dirac masses strictly equal to the up-type quark masses and fitting to the neutrino oscillation parameters generally yields strongly hierarchical right-handed neutrino masses ($M_1 \ll M_2 \ll M_3$), with $M_1 \sim 10^5$ GeV. The lepton asymmetry in this case is too small by several orders of magnitude. However, it is plausible that there are large radiative corrections to the

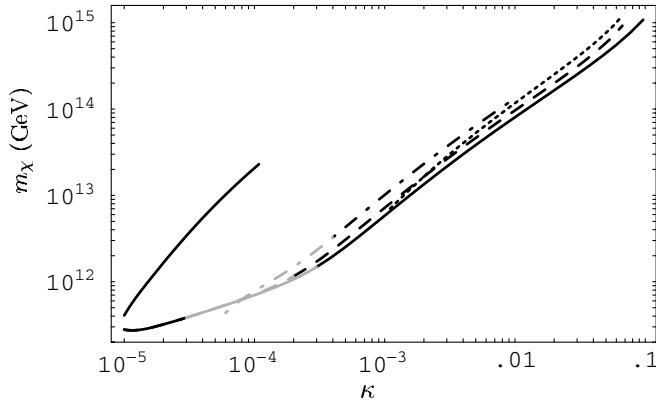


FIG. 5. The inflaton mass m_χ vs the allowed range of κ ($n_s < 1.04$), for SUSY hybrid inflation with $\mathcal{N} = 1$ (solid line), with $\mathcal{N} = 2$ (dashed line), and for shifted hybrid inflation (dot-dashed line for $M_S = m_p$, dotted line for $M_S = 5 \times 10^{17}$ GeV). The gray segments denote the range of κ for which the change in $\arg S$ is significant.

first two family Dirac masses, so that M_1 remains heavy compared to T_r .

A reasonable mass pattern is therefore $M_1 < M_2 \ll M_3$, which can result from either the dimensionless couplings γ_{ij} or additional symmetries (see, e.g., [25]). The dominant contribution to the lepton asymmetry is still from the decays with N_3 in the loop, as long as the first two family right-handed neutrinos are not quasidegenerate. Under these assumptions, the lepton asymmetry is given by [26]

$$\frac{n_L}{s} \lesssim 3 \times 10^{-10} \frac{T_r}{m_\chi} \left(\frac{M_i}{10^6 \text{ GeV}} \right) \left(\frac{m_{\nu 3}}{0.05 \text{ eV}} \right), \quad (14)$$

where M_i denotes the mass of the heaviest right-handed neutrino the inflaton can decay into. The decay rate $\Gamma_\chi = (1/8\pi)(M_i^2/M^2)m_\chi$ [7], and the reheat temperature T_r is given by

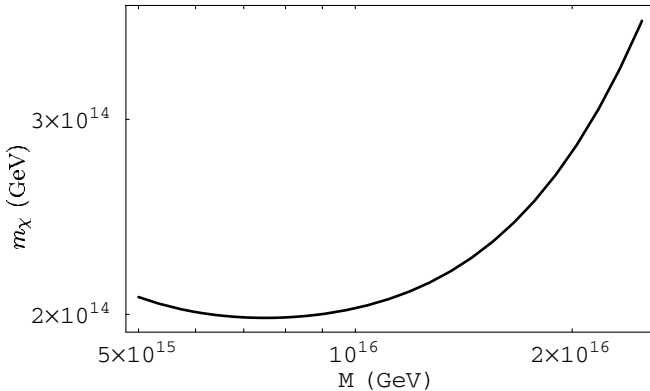


FIG. 6. The inflaton mass m_χ vs the symmetry breaking scale M for smooth hybrid inflation.

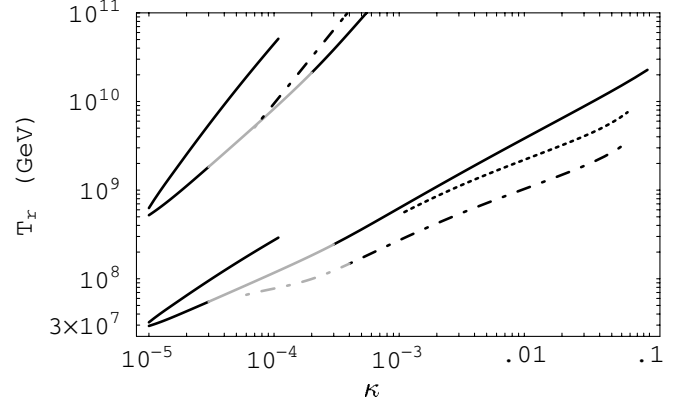


FIG. 7. The lower bound on the reheat temperature T_r vs the allowed range of κ ($n_s < 1.04$), for SUSY hybrid inflation with $\mathcal{N} = 1$ (solid line) and for shifted hybrid inflation (dot-dashed line for $M_S = m_p$, dotted line for $M_S = 5 \times 10^{17}$ GeV). The segments in the top left part of the figure correspond to the bounds in the presence of a $\lambda S h^2$ coupling. The gray segments denote the range of κ for which the change in $\arg S$ is significant.

$$T_r = \left(\frac{90}{\pi^2 g_*} \right)^{1/4} (\Gamma_\chi m_p)^{1/2} \simeq \frac{1}{10} \frac{(m_p m_\chi)^{1/2}}{M} M_i. \quad (15)$$

(We have ignored the effect of preheating in hybrid inflation [27], which does not seem to change the perturbative estimate for T_r significantly [28].) From the experimental value of the baryon to photon ratio $\eta_B \simeq 6.1 \times 10^{-10}$ [14], the required lepton asymmetry is found to be $n_L/s \simeq 2.5 \times 10^{-10}$ [29]. Using this value, along with Eqs. (14), (15), we can express T_r in terms of the symmetry breaking scale M and the inflaton mass m_χ :

$$T_r \gtrsim 1.9 \times 10^7 \text{ GeV} \left(\frac{10^{16} \text{ GeV}}{M} \right)^{1/2} \left(\frac{m_\chi}{10^{11} \text{ GeV}} \right)^{3/4} \times \left(\frac{0.05 \text{ eV}}{m_{\nu 3}} \right)^{1/2}. \quad (16)$$

Here m_χ is given by $\sqrt{2}\kappa M$, $\sqrt{2}\kappa M\sqrt{1-4\xi}$ and $2\sqrt{2}v^2/M$, respectively, for hybrid, shifted hybrid and smooth hybrid inflation. The value of m_χ is shown in Figs. 5 and 6. We show the lower bound on T_r calculated using this equation (taking $m_{\nu 3} = 0.05$ eV) in Figs. 7 and 8.

Equation (15) also yields the result that the heaviest right-handed neutrino the inflaton can decay into is about 300 (5) times heavier than T_r , for hybrid inflation with $\kappa = 10^{-5}$ (10^{-2}). For shifted hybrid inflation, this ratio does not depend on κ as strongly and is $\sim 10^2$ [9]. This is consistent with ignoring washout effects as long as the lightest right-handed neutrino mass M_1 is also $\gg T_r$.

Both the gravitino constraint and the constraint $M_1 \gg T_r$ favor smaller values of κ for hybrid inflation, with $T_r \gtrsim 3 \times 10^7$ GeV for $\kappa \sim 10^{-5}$. Similarly, the gravitino constraint favors κ values as small as the inflationary trajectory

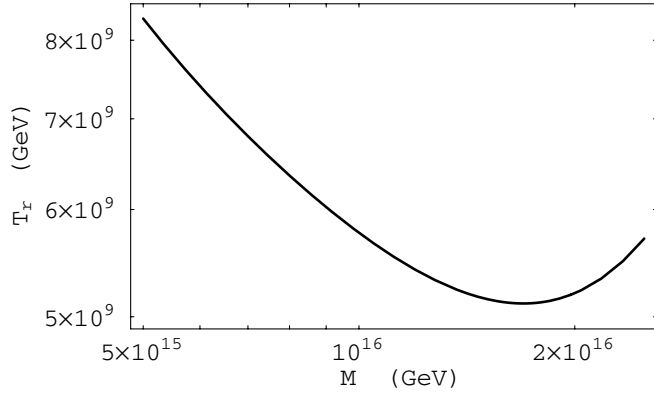


FIG. 8. The lower bound on the reheat temperature T_r vs the symmetry breaking scale M for smooth hybrid inflation.

allows for shifted hybrid inflation, and $T_r \gtrsim 10^8$ GeV for $M_S = m_p$. Smooth hybrid inflation is relatively disfavored since $T_r \gtrsim 5 \times 10^9$ and $M_2/T_r \approx 3(9)$ for $M = 5 \times 10^{15}$ GeV (2×10^{16} GeV).²

There are ways to evade these bounds on T_r . Having quasidegenerate neutrinos increases the lepton asymmetry per neutrino decay ϵ [30] and thus allows lower values of T_r corresponding to lighter right-handed neutrinos. Provided that the neutrino mass splittings are comparable to their decay widths, ϵ can be as large as $1/2$ [31]. The lepton asymmetry in this case is of order T_r/m_χ , where $m_\chi \sim 10^{11}$ GeV for $\kappa \sim 10^{-5}$, and sufficient lepton asymmetry can be generated with T_r close to the electroweak scale. For other scenarios that yield T_r of order 10^6 GeV in hybrid inflation, see [8] (without a $B - L$ symmetry) and [32].

We end this section with some remarks on the μ problem and the relationship to T_r in the present context. The MSSM μ problem can naturally be resolved in SUSY hybrid inflation models in the presence of the term λSh^2 in the superpotential, where h contains the two Higgs doublets [33]. [The “bare” term h^2 is not allowed by the $U(1)$ R-symmetry.] After inflation the vev of S generates a μ term with $\mu = \lambda \langle S \rangle = -m_{3/2} \lambda / \kappa$, where $\lambda > \kappa$ is re-

TABLE I. Lower bounds on the reheat temperature (GeV).

	Without λSh^2	With λSh^2
SUSY hybrid inflation	3×10^7	5×10^8
Shifted hybrid inflation	7×10^7	5×10^9
Smooth hybrid inflation	5×10^9	$\gtrsim 10^{12}$

quired for the required vacuum. The inflaton in this case predominantly decays into Higgses (and Higgsinos) with $\Gamma_h = (1/16\pi)\lambda^2 m_\chi$. As a consequence, the presence of this term significantly increases the reheat temperature T_r . Following [34], we calculate T_r for the best case scenario $\lambda = \kappa$. We find a lower bound on T_r of 5×10^8 GeV in hybrid inflation, see Fig. 7. $T_r \gtrsim 5 \times 10^9$ GeV for shifted hybrid inflation with $M_S = m_p$,³ and $T_r \gtrsim 10^{12}$ GeV for smooth hybrid inflation. An alternative resolution of the μ problem in these models which has no impact on T_r invokes an axion symmetry [3,35].

IV. CONCLUSION

Supersymmetric hybrid inflation models, through their connection to the grand unification scale, provide a compelling framework for the understanding of the early universe. Such models can also meet the gravitino and baryogenesis constraints through nonthermal leptogenesis via inflaton decay. SUGRA corrections and previously ignored soft SUSY violating terms in the inflationary potential lead to lower bounds on κ and therefore, assuming hierarchical right-handed neutrinos, on the reheat temperature T_r . The lower bounds on T_r are summarized in Table I.

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²A new inflation model related to smooth hybrid inflation is discussed in [25], where the energy scale of inflation ν is lower and consequently lower reheat temperatures are allowed.

³We take $\lambda/\kappa = 2/(1/4\xi - 1)$ for shifted hybrid inflation. Some scalars belonging to the inflaton sector acquire negative mass² if λ is smaller. $\kappa \sim 10^{-4}$ corresponds to $\lambda/\kappa \approx 3$.

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