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Constraints on Supergravity Chaotic Inflationary Models

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Abstract

We discuss, in the context of $N = 1$ hidden sector supergravity models, constraints on the parameters of a polynomial superpotential resulting from existing bounds on the reheating temperature and on the amplitude of the primordial energy density fluctuations as inferred from COBE. We present a specific two-parameter chaotic inflationary model which satisfies these constraints and discuss a possible scenario for adequate baryon asymmetry generation.

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Recently, a great deal of attention has been devoted to possible implementations of the inflationary scenario in supergravity/superstring models. Besides ensuring that there is sufficient inflation to solve the initial condition problems of the standard cosmological model, these models have to explain observational limits such as the magnitude of the energy density perturbations required to explain the anisotropies in the Cosmic Microwave Background radiation observed by COBE [1]. Furthermore, in supergravity cosmological models, the reheating temperature should not exceed $T_{RH} \lesssim 2.5 \times 10^8 (100\ \mathrm{GeV}/m_{3/2})$ GeV [2], not to generate an abundance of gravitinos which would photo-dissociate light elements produced in primordial nucleosynthesis. In the context of superstring cosmology, inflationary models have to face further problems such as the fate of the dilaton and moduli fields and the so-called post-modern Polonyi problem [3]. Although many of these issues can be addressed in a simple chaotic model where the dilaton plays the role of the inflaton and its potential is dominated by quadratic and/or quartic self-couplings [4], this is not the case for the dilaton potentials generated by the supersymmetry-breaking mechanisms currently preferred in superstring-based models, i.e. gaugino condensation; moreover, these potentials appear not to be sufficiently flat to allow inflation to occur [5].

In this letter, we study constraints on $N = 1$ supergravity chaotic inflationary models arising from the superstring, resulting from the abovementioned cosmological bounds. Chaotic inflationary models [6] stand out as the most natural ones in what the initial conditions for the onset of inflation are concerned, particularly in the context of supergravity and superstring theories, where the natural scale for fields is the Planck scale. Realizations of chaotic inflation in minimal and in $SU(1,1)$ $N = 1$ supergravity theories have been studied in [7]. We discuss a specific model, with a two-scale chaotic inflationary sector, which could originate from the existence of two gaugino condensation and/or gauge symmetry breaking scales, that can accommodate in a satisfactory way the bounds on the reheating temperature and energy density fluctuations. It follows from our analysis that, as first pointed out in Ref. [8], a chaotic inflationary model requires more than one scale to reproduce the abovementioned constraints; this is essentially due to the fact that in chaotic models the slow roll-over period occurs around the Planck scale and not, as in Refs. [8, 9], some orders of magnitude below.

We shall assume that the inflaton is the scalar component of a gauge singlet superfield, Φ, in the hidden sector of the theory. We start by splitting the superpotential in a supersymmetry-breaking, a gauge and an inflationary part, as suggested in [8, 9]:

$$W = P + G + I. \tag{1}$$

The scalar potential for the inflaton field is obtained from the superpotential $I(\Phi)$ as

$$V(\phi) = \exp(-|\phi|^2/M^2)\,\left(|\partial I/\partial\Phi + M^{-2}\Phi^* I|^2 - 3M^{-2}|I|^2\right)\Big|_{\Phi=\phi}. \tag{2}$$

with $M = M_P/\sqrt{8\pi}$, where M_P is the Planck mass. Requiring the cosmological constant to vanish and that supersymmetry remains unbroken at the minimum of the potential, $\Phi = \Phi_o$, leads to the following constraints on the superpotential:

$$I(\Phi_o) = \frac{\partial I}{\partial \Phi}(\Phi_o) = 0. \tag{3}$$

Consider the most general polynomial superpotential

$$I(\Phi) = \sum_{n=0} \frac{a_n}{M^{n-2}} \Phi^n, \tag{4}$$

where the a_n are mass parameters. Dropping the linear and the non-renormalizable terms ($n > 3$) in (4), the latter leading to too large tensor perturbations of the microwave background [10], we are left with

$$I(\Phi) = I_o + a\;\Phi^2 + \frac{b}{M}\;\Phi^3\;, \tag{5}$$

where I_o is a constant and a and b are positive. The conditions (3), applied to the superpotential of Eq. (5), give two solutions for ϕ_o, I_o:

$$\begin{aligned} (\phi_o,\ I_o) &= (0,\ 0); &\quad (6)\\ &= (-2aM/3b,\ -4a^3M^2/27b^2). &\quad (7)\end{aligned}$$

We shall first consider the case $\phi_o = I_o = 0$. The inflaton potential (along the real ϕ direction) is then given by

$$\begin{aligned} V(\phi) &= M^2 \exp\left(-\phi^2/M^2\right)\left[4\;a^2\left(\frac{\phi}{M}\right)^2 + 12\;a\;b\left(\frac{\phi}{M}\right)^3 + (a^2+9\;b^2)\left(\frac{\phi}{M}\right)^4\right.\\ &\left. + 4\;a\;b\left(\frac{\phi}{M}\right)^5 + (a^2+3\;b^2)\left(\frac{\phi}{M}\right)^6 + 2\;a\;b\left(\frac{\phi}{M}\right)^7 + b^2\left(\frac{\phi}{M}\right)^8\right]. \end{aligned} \tag{8}$$

In the chaotic inflationary scenario, the scalar field starts rolling towards its minimum at the origin from an initially large value beyond the Planck scale. During this process, the domains of the Universe filled with a sufficiently homogeneous ϕ field expand according to the Friedmann equation

$$H^2 = \frac{1}{3M^2}\left(\frac{1}{2}\dot{\phi}^2 + V\right), \tag{9}$$

and the ϕ field evolves according to

$$\ddot{\phi} + 3H\dot{\phi} + \frac{dV}{d\phi} = 0. \tag{10}$$

In the region $|\phi| >$ few M, the field rolls down very slowly and the terms $\dot{\phi}^2$ in (9) and $\ddot{\phi}$ in (10) can be neglected. In this region, the potential $V(\phi)$ can be approximated by

$$V(\phi) \approx \frac{b^2}{M^6}\phi^8 \exp\left(-\phi^2/M^2\right). \tag{11}$$

The total number of e-folds of inflation is given by

$$N \equiv \ln\frac{a(\phi_e)}{a(\phi_i)} = -\frac{1}{M^2}\int_{\phi_i}^{\phi_e}\frac{V}{V'}d\phi \approx \frac{\pi}{2M_P^2}\left(\phi_i^2 - \phi_e^2\right). \tag{12}$$

Hence, it is required that $\phi_i \gtrsim 6.4 M_P$, for $\phi_e \approx M_P$, to get $N \gtrsim 65$.

After inflation, the field ϕ begins to oscillate about its minimum, thus reheating the Universe. At minimum, the inflaton field has a mass

$$m_\phi = 2\sqrt{2}\ a. \tag{13}$$

Since the inflaton is hidden from the other sectors of the theory, it couples to lighter fields with strength $\sim a/M$, leading to a decay width

$$\Gamma_\phi \sim \frac{m_\phi}{(2\pi)^3}\left(\frac{a}{M}\right)^3, \tag{14}$$

and a reheating temperature

$$T_{RH} \approx \left(\frac{30}{\pi^2 g_{RH}}\right)^{1/4}\sqrt{M\Gamma} \sim \frac{2}{\pi^2}\left(\sqrt{\frac{15}{g_{RH}}\frac{a^3}{M}}\right)^{1/2}, \tag{15}$$

where g_{RH} is the number of degrees of freedom at T_{RH}. Notice that, as $\Gamma_\phi << m_\phi$, parametric resonance effects [11] are not important in this case.

As mentioned above, a severe upper bound on T_{RH} comes from the requirement that sufficiently few gravitinos are regenerated in the post-inflationary reheating epoch. Indeed, once regenerated beyond a certain density, stable thermal gravitinos would dominate the energy density of the Universe or, if they decay, have undesirable effects on nucleosynthesis and light element photo-dissociation and lead to distortions in the microwave background. This implies the following bounds[3] [2, 14]:

$$T_{RH} \lesssim 2 \times 10^9,\ 6 \times 10^9 \text{ GeV} \quad \text{for} \quad m_{3/2} = 1,\ 10 \text{ TeV}. \tag{16}$$

For our model, demanding that T_{RH} be less than 2×10^9 GeV then leads, for $g_{RH} \approx 150$, to a bound on parameter a as

$$\frac{a}{M} \lesssim 3.7 \times 10^{-6}. \tag{17}$$

Further constraints on the parameters of the superpotential can be derived from the spectrum of adiabatic density fluctuations, which is given, in terms of the potential, by [15]:

$$\delta_H \equiv \left(\frac{\delta\rho}{\rho}\right)_H = \frac{1}{5\sqrt{3}\pi M^3} \frac{V_\star^{3/2}}{V_\star'}, \tag{18}$$

where the subscript $\star$ indicates that the right-hand side should be evaluated as the comoving scale k equals the Hubble radius ($k = aH$) during inflation. For a sufficiently flat spectrum and no significant generation of long wavelength gravitational waves, the central value of the 10° anisotropy observed by COBE is reproduced, provided [1]

$$\delta_H \approx 2.3 \times 10^{-5}. \tag{19}$$

Combining (18) and (19), we obtain for our model

$$\delta_H \approx \frac{1}{5\sqrt{3e}\ \pi\ M} \frac{(6\ a^2 + 18\ a\ b + 13\ b^2)^{3/2}}{(3\ a^2 + 17\ a\ b + 18\ b^2)}, \tag{20}$$

which, using the constraint on a derived above, Eq. (17), implies, in turn, a bound on parameter b

$$\frac{b}{M} \lesssim 7.8 \times 10^{-4}. \tag{21}$$

[3]Fischler [12] suggested that heat-bath effects might greatly enhance the gravitino regeneration rate at high temperature and thereby lower the bound on T_{RH}, a claim that has since been questioned [13].

For the other solution of Eq. (3), i.e. $\phi_o = -2aM/3b$, $I_o = -4a^3/27b^2$, we obtain the same bounds on a and b, as expected.

Of course, our results would be modified if, instead of (16), there were stricter bounds on the reheating temperature [12]. For instance, for $T_{RH} \lesssim 10^6$ GeV, we obtain:

$$\frac{a}{M} \lesssim 2.3 \times 10^{-8} \ , \ \frac{b}{M} \lesssim 7.8 \times 10^{-4}. \tag{22}$$

A realistic scenario for baryogenesis can be built considering the decay of the inflaton into the matter field states in the gauge sector of the superpotential (1). As the coupling between the inflaton and these states is only gravitational, the former will decay into the heaviest states available [9], which will then generate the baryon asymmetry through decays into quarks and leptons. The baryon-antibaryon number density will then be given essentially in terms of the asymmetry following from inflaton decay:

$$n_{B-\bar{B}} \ \approx \ n_\phi \, \delta B \ , \tag{23}$$

where δB is the baryon asymmetry generated per decay. The photon number density can be given as a function of the inflaton density and the reheating temperature

$$n_\gamma \ \approx \ \frac{\rho_\phi}{T_{RH}} \ \approx \ \frac{n_\phi m_\phi}{T_{RH}} \ , \tag{24}$$

so that, from (13), (16) and (17), the asymmetry can be expressed as

$$\xi \equiv \frac{n_{B-\bar{B}}}{n_\gamma} \ \approx \ \frac{T_{RH}}{m_\phi} \, \delta B \sim 10^{-4} \ \delta B \ , \tag{25}$$

which allows us to obtain the observed value, $\xi \sim 10^{-10}$, provided δB has a suitable value. Although we shall not try to specify here how, in a concrete particle physics model, the required value for δB could be produced radiatively, we stress that the asymmetry can be created even though the reheating temperature is as low as or lower than the bound (16). We also point out that, in models such as the ones discussed in Ref. [4] (see also [16]), the inflaton (the dilaton, in that instance) is directly coupled to a GUT Higgs field and the mechanism discussed above can be easily implemented. This GUT Higgs field can be endowed with a suitable potential that may allow a subsequent period of inflation.

Of course, other scenarios could be envisaged to generate the baryon asymmetry, which can be completely or fairly independent of inflaton decay, depending on whether or not this decay dilutes the generated baryon asymmetry (see e.g. the second reference in [4]). An example is the Affleck-Dine mechanism [17], recently implemented in the context of supergravity string-inspired models [18]. The main feature invoked in these models is that, during inflation, supersymmetry-breaking soft terms, with mass terms of the order of the Hubble parameter, are naturally induced [19]. Since the cosmology of string theories has problems associated with the fate of the moduli fields, an additional period of late inflation seems to be a rather natural way to avoid the problems associated with the presence of these fields [20]. The possibility that this late period of inflation is related with baryon asymmetry generation itself is certainly very appealing. In the context of our model, an Affleck-Dine baryogenesis scenario like the one in [18] can also be constructed. Non-renormalizable terms, together with soft supersymmetry-breaking terms arising from a second period of GUT inflation, give rise, along some direction in the space of scalar fields that carry baryon and lepton number (χ), such as squarks and sleptons, to the potential [18, 19]:

$$V(\chi) \approx c\, H^2|\chi|^2 + a\, \lambda \frac{H|\chi|^n}{nM^{n-3}} + \lambda^2 \frac{|\chi|^{2n-2}}{M^{2n-6}} \quad , \tag{26}$$

where a and c are O(1) constants – "a terms" are important for B and L violation – and λ is a coupling constant. This potential admits a non-trivial minimum $|\chi_0| \approx \left(\frac{-c}{n-1}\right)^{1/2} \frac{H}{\lambda}\, M^{n-3}$, for $c < 0$. After the second period of inflation, when $H \approx m_{3/2}$, the field oscillates around χ_0 and a baryon asymmetry such as (25) is generated, with δB given essentially by $(\chi_0/M_P)^2$ [18].

Finally, we comment on possible origins for the (two) scales of our model. Let us first consider the possibility that these scales are induced by gauge symmetry breaking. In fact, as suggested in Ref. [8], once gauge non-singlet fields, Ψ, acquire a v.e.v. along a D-flat direction, thereby breaking the gauge symmetry, a v.e.v. for the massive gauge singlet fields coupled to them will be induced; these v.e.v.'s then feed through to the inflationary sector via couplings between the latter fields and the inflaton, leading to a superpotential of the form [8]

$$I(\Phi) = a\, M^2\, f\left(\frac{\Phi}{M}\right), \tag{27}$$

where $a = \langle\Psi\rangle^2\langle\bar{\Psi}\rangle^2/M^3$ and $f(x)$ is a polynomial function. For $\langle\Psi\rangle \sim 10^{16}$ GeV, we obtain $a \sim 10^{10}$ GeV. Of course, for our model, at least two operators coupling Ψ and Φ fields would be required, whose form is determined by e.g. discrete symmetries, which arise naturally in the context of string-inspired phenomenological models as a consequence of the symmetries of the compactification manifold [21]. In fact, one hopes that these and/or other symmetries present in the fundamental theory may also explain the absence of higher-order non-renormalisable terms, which are not small in chaotic inflationary models, where $\Phi \geq M_P$. Other possibilities for the origin of these scales is that there are two stages of gauge symmetry breaking or that they arise from the supersymmetry-breaking sector, through the gaugino condensation of two distinct subgroups of the hidden group E_8.

Let us now summarize our results. We have shown that, in order to satisfy COBE data and to keep the reheating temperature sufficiently low not to regenerate an excessive abundance of gravitinos, chaotic inflationary models require at least two independent scales in the superpotential. These scales, which can be related with the scales of gaugino condensation and/or gauge symmetry breaking, are significantly below the Planck scale. We have also analysed how inflaton decay can potentially explain the observed baryon asymmetry of the Universe.

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