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SUPERSYMMETRIC TWO-STAGE INFLATION[1]

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We investigate a supersymmetric extension of the hybrid potential with the mass-term replaced by a self-coupling quartic one. This gives a scenario with two disconnected stages of inflation. The first is an usual chaotic inflation, while the second is a vacuum-dominated one. However, in this model the latter stage requires a fine-tuning of the self-coupling constant in order to occur and to generate the density perturbations.

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1 Introduction

Inflation [1] was initially invoked as a solution to the longstanding problems of the standard hot big-bang model as the horizon and flatness problems. These are related to the necessity of imposing severe initial conditions. Some other problems like the monopole one have been also solved in the context of inflation. Inflation also provides a mechanism for the generation of primordial cosmological perturbations, which are responsible for the observed temperature anisotropies in the cosmic microwave background radiation (cmbr) and the large-scale structure of the Universe. The initial version of Guth [2] was based on a first order phase transition, but in improved models [3, 4] the inflationary phase is associated with a second order phase transition accompanied by a period of slow-rolling of the inflaton (the scalar field driving inflation). This was the basic mechanism for all the inflationary models proposed so far.

Succesful implementation of the inflationary picture requires a long enough period of rapid expansion and a correct order of magnitude for primordial perturbations. In the simplest version of single-field inflation (the chaotic one [4]), the corresponding constraints on the potential of the inflaton (the scalar field driving inflation) are unrealistic from a particle physics point of view. Indeed, coupling parameters must be fine tuned to very small values, while the inflaton must be of the order of the Planck mass during inflation.

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These problems can be avoided in the so-called hybrid inflation scenario, recently proposed by Linde [5], where more than one field is relevant. This is one of the most attractive model of inflation at present. It shares the best feature of the slow-rolling and the first order phase transition models. Indeed, during inflation, the non inflaton field is trapped in a false vacuum state and the Universe is dominated by the false vacuum energy. Inflation ends with a first order phase transition when the inflaton field reaches some critical value.

Furthermore, in nonsupersymmetric theories, a flat potential dos not emerge naturally (except in the case of Goldstone boson). Contrarily, it is well known that supersymmetric theories often admit non compact flat vacuum directions [6]. These directions remain flat to all order of perturbation theory, in the unbroken SUSY limit, due to the non-renormalisation theorem, but they gain a very small curvature induced by SUSY breaking which is very useful for inflation. The supersymmetric flat directions are then good candidates for inflation.

Hybrid inflation is naturally realised in supersymmetric theories and its relevance to SUSY has been extensively investigated [7–9], Many alternatives have been proposed [10–12]. In supersymmetric hybrid models the inflaton which is usually a scalar singlet is coupled to a Higgs superfield that is charged under some gauge group G. At the end of inflation the Higgs field receives a non vanishing vev, G is then spontaneously broken.

The non-zero vacuum energy density during inflation can either be due to the vev of an $F-$term or that of a $D-$term [13]. The scalar potential has two minima : a local minimum of value of the inflaton S greater than some critical value S_c (with a vanishing Higgs field), and a global supersymmetric minimum at $S = 0$ and a non-zero Higgs vev. The two fields are assumed to have chaotic initial conditions. While $S \gg S_c$, the Higgs field rapidly setlles down to the local minimum. The SUSY is broken since the Universe is dominated by a non vanishing vacuum energy. The slow-rolling conditions are satisfied and inflation takes place until $S = S_c$. When S falls below S_c the Higgs field starts to acquire a non-zero vev. both fields then oscillate before stabilising at the global supersymmetric minimum. However, the initial supersymmetric realization of hybrid inflation [7] gives a scalar potential without a slope in the inflaton direction which means that inflation could not end in such a model.

In the present paper we propose a solution to this problem based on an extension of the usual superpotential for supersymmetric hybrid inflation. The corresponding model gives rise to two disconnected stages of inflation. The first one corresponds to a simple chaotic inflation scenario, while the second is a vacuum-dominated inflation. The remainder of the paper is organized as follows: in section 2 we give a brief review of the initial version of the hybrid inflation in both supersymmetric and nonsupersymmetric cases. In section 3 the two-stage inflation model is presented. Section 4 is devoted to the discussion and conclusions.

2 Hybrid Inflation

The hybrid inflation is a multifield model which was introduced initially to overcome one of the difficulties of the usual inflationary models based on the slow-rolling potentials with a single field, namely how to make the potential very flat and to give rise to a rapid end of inflation and a sufficient reheating. By using two fields instead of one, this problem was beautifully solved. One field provides the vacuum energy which drives inflation, while the second acts as the inflaton.

The model is based on a scalar potential of the form [5, 8]

$$V(\varphi,\psi)=\kappa^2\Big(\mu^2-\frac{\varphi^2}{4}\Big)^2+g^2\frac{\varphi^2\psi^2}{4}+m^2\frac{\psi^2}{2} \tag{1}$$

where κ, g are dimensionless coupling constants and μ,m are mass parameters. The global minimum is at:

$$\varphi=\pm 2\mu \tag{2}$$

$$\psi=0$$

There is also a false vacuum at $\varphi=0$ with $V(\varphi=0,\psi)=\kappa^2\mu^4+m^2\frac{\psi^2}{2}$. In this minimum the mass squared of the φ field is:

$$m^2_\varphi=-\kappa^2\mu^2+\frac{1}{2}g^2\psi^2 \tag{3}$$

This mass is positive as long as $\psi>\psi_c$ where ψ_c is given by:

$$\psi_c=\sqrt{2}\frac{\kappa\mu}{g} \tag{4}$$

In the absence of the mass term of ψ in eq. (1) we obtain an exact flat direction at $\varphi=0$ and $\psi>\psi_c$ along which the inflation may occur. But when ψ drops below ψ_c, m^2_φ becomes negative and the symmetry is broken. Then the field φ will rapidly move toward its true vacuum value (2) leading to a waterfall end of inflation. When we reintroduce the mass term, the flat direction acquires a non zero slope which drives the inflaton field ψ to its value at the end of the slow-rolling phase, otherwise the inflation will never end.

In the above discussion the end of inflation was identified with $\psi=\psi_c$. However, this can happen well before this instant, when the slow-rolling conditions [14] cease to be valid as we will see later.

The idea of hybrid inflation was also studied in the context of supersymmetric theories by many authors [7–9]. The model is realized by the superpotential [7]

$$W=\kappa S\big(-\mu^2+\overline{\phi}\phi\big) \tag{5}$$

where ϕ,$\overline{\phi}$ is a conjugate pair of chiral superfields which belong to a non-trivial representation of a group G, and S is a gauge singlet superfield. This superpotential is the most general renormalizable one consistent with a $U(1)_R$ symmetry (under which $W\longrightarrow e^{i\theta}W$, $S\longrightarrow e^{i\theta}S$ and $\overline{\phi}\phi\longrightarrow\overline{\phi}\phi$). Using the expression [15]

$$V\;=\;\sum_i\left|\frac{\partial W}{\partial\phi_i}\right|^2\;+\;D-\text{Term} \tag{6}$$

the scalar potential is given by

$$V\;=\;\kappa^2\big|\mu^2\;-\;\overline{\phi}\phi\big|^2\;+\;\kappa^2\big|S\big|^2\big(\big|\phi\big|^2\;+\;\big|\overline{\phi}\big|^2\big)\;+\;D-\text{Term} \tag{7}$$

We can take the $D-$term vanishing if we restrict ourselves to the D-flat direction $\phi^* = \phi$ which contains the supersymmetric minima.

If we set:

$$S \equiv \frac{\psi}{\sqrt{2}} \tag{8}$$

$$\overline{\phi} = \phi \equiv \frac{\varphi}{2} \tag{9}$$

where ψ and φ are normalized real scalar fields. The scalar potential derived from the superpotential(5) then takes the form (1) with $\kappa = g$ and $m = 0$.

The potential for hybrid inflation is therfore obtainable from SUSY but without the mass-term of ψ which is necessary for inflation since, as we have noted above, it gives the slope of the valley of minima.

One way to obtain a useful flat direction for inflation (with a nonvanishing slope) is to generate a potential for the S field through radiative corrections [7]. This is possible when the supersymmetry is broken, in the exact supersymmetry limit the non renormalization theorem protects the flat direction against such corrections.

When we consider the radiative corrections, the one loop effective potential for $\psi >> \psi_c$ is then given by [7]

$$V_{\text{eff}}(\psi) = \kappa^2\mu^4\Big[1 + \frac{\kappa^2}{32\pi^2}\Big(\ln\frac{\kappa^2\psi^2}{2\Lambda^2} + \frac{3}{2}\Big)\Big] \tag{10}$$

which provides a small slope in the ψ direction that will drive ψ to its value at the end of the slow-rolling phase.

Another way is to modify the superpotential of Eq.(5) by using the first nonrenormalizable contribution instead of renormalizable trilinear coupling. The new superpotential is then [21–23]

$$W \;=\; S\bigg(-\mu \;+\; \frac{\langle\phi\overline{\phi}\rangle^2}{M^2}\bigg) \tag{11}$$

where M is a mass scale of the order of the compactification scale $M_c \sim 10^{18}$ GeV which controls the nonrenormalizable terms in the superpotential of the theory. The scalar potential obtained from W of Eq.(11) is;

$$V(\chi\,,\,\sigma) \;=\; \bigg(\mu^2 \;-\; \frac{\chi^2}{16M^2}\bigg)^2 \;+\; \frac{\chi^2\sigma^2}{16M^4} \tag{12}$$

Although this potential is very similar to the one of the Eq.(1) with $m = 0$, it possesses some crucial differences. The flat direction at $\chi = 0$ with $V(\chi = 0, \sigma) = \mu^4$ is now a local maximum in the χ direction for all values of σ, but two valleys of local minima developp on both sides and close to this flat direction at $\chi^2 \simeq 4\mu^2M^2/3\sigma^2$ ($\sigma^2 \gg \mu M$) with $V \simeq \mu^4(1 - 2\mu^2M^2/27\sigma^4)$. These valleys have an inbuilt slope and thus they can in principle be used for inflation [21, 23]. This is called a smooth hybrid inflation model.

3 Two-Stage Inflation

3.1 The model

Our main motivation is to contribute by a new alternative to the resolution of the above mentioned problem of the vanishing mass of the inflaton by a modification of the superpotential of Eq.(5). Our model is based on a superpotential of the form:

$$W = \kappa S\big(-\mu^2 + \overline{\phi}\phi\big) + \frac{\lambda_s}{3} S^3 \tag{13}$$

This is a simple extension of the superpotential (5) which involves a renormalizable self-coupling of the singlet scalar superfield[3].

The scalar potential derived from the above superpotential is then given by:

$$V(\psi,\varphi) = \kappa^2\Big(\mu^2 - \frac{\varphi^2}{4}\Big)^2 + \big(g^2 + \kappa\lambda_s\big)\frac{\varphi^2\psi^2}{2} + \lambda\frac{\psi^4}{4} + \kappa\lambda_s\mu^2\psi^2 \tag{14}$$

where $\lambda = \lambda_s^2$, and where the fields definition is the same as in (8) and (9).

As we shall see later, the typical order of magnituede of the coupling constant λ is $O(10^{-14})$, and with the choice $g, \kappa \sim O(10^{-1})$ we can take:

$$\Big(\frac{g^2}{2} + \kappa\lambda_s\Big) \sim \frac{g^2}{2} \tag{15}$$

Since during the inflationary phase we are interested in the large values of ψ($\psi >> \psi_c$), the mass-term $\kappa\lambda_s\mu^2\psi^2$ can be neglected with respect to the ψ^4-term. The mass-term can be considered in a class of hybrid inflation called inverted hybrid inflation [16] where the inflaton rolls away from the origin.

The effective potential can then be approximated by:

$$V(\psi,\varphi) = \kappa^2(\mu^2 - \frac{\varphi^2}{4})^2 + g^2\frac{\varphi^2\psi^2}{4} + \lambda\frac{\psi^4}{4} \tag{16}$$

in these conditions.

This scalar potentiel has the same form as the one of the initial Linde's model (1) but the mass-term is now replaced by a $\lambda\psi^4$-term.

3.2 Two-stage inflation

Since the $\lambda\psi^4$-term and the $m^2\psi^2$-term are the two familiar terms in the chaotic one field inflationary models we can expect the above potential to describe the same dynamics as the model (1). However, as we shall see later, this model gives rise to two disconnected stages of inflation. A similar result has been found by the authors of the reference [17] but that model was based on two massive scalar fields with no direct interaction betweeen them, and each stage of inflation is drived by one of them. In such models the spectrum of density perturbations may have a very rich and nontrivial structure.

[3] Although some authors usually invoke a continuous $R-$symmetry that forbids such a term, for phenomenological reasons this symmetry can not be used to exclud this trilinear term [24]

As for Linde's model the above potential presents a global minimum at $\psi = 0$ and $\varphi = \pm 2\mu$ and a flat direction for $(\psi > \psi_c, \varphi = 0)$, along which the potential is:

$$V(\varphi = 0, \psi) = \kappa^2 \mu^4 + \lambda \frac{\psi^4}{4} \tag{17}$$

At large ψ the potential (17) is dominated by the second term ($\lambda \frac{\psi^4}{4} >> \kappa^2 \mu^4$) and the model reduces to a simple chaotic one [4]. In the chaotic inflation we only impose, as initial condition for the inflaton field:

$$\lambda \frac{\psi^4}{4} \lesssim m_p^4 \tag{18}$$

in order to insure the validity of the classical description of the Universe, this implies that $\sqrt{\kappa}\mu \ll m_p$, which is realizable by the choice $\mu \sim 10^{15}$ GeV (this choice is dictated by some particle physics considerations, in particular to give a connection with the succession of the phase transitions in the early Universe), the potentiel is dominated by the $\psi-$term. The equation of motion for the scalar field in a flat isotropic Friedmann Universe is:

$$\ddot{\psi} + 3H\dot{\psi} = -\frac{\partial V}{\partial \psi} \tag{19}$$

where $H^2 \simeq \frac{8\pi}{3m_p^2} V$. Thus, if we assume the existence of some region of the Universe of size $l \sim H^{-1}$ in which the field ψ is sufficiently homogeneous and isotropic with the initial condition (18), this region undergoes a period of chaotic inflation if we can neglect the first term in the Eq.(19).

The inflation takes place as long as we have [14]

$$\epsilon \equiv \frac{m_p^2}{16\pi} \left(\frac{V'(\psi)}{V(\psi)} \right)^2 \ll 1 \tag{20}$$

$$\eta = \frac{m_p^2}{8\pi} \frac{V''(\psi)}{V(\psi)} \ll 1 \tag{21}$$

Since our potential is reduced to a quartic one during this stage of inflation we can write $\epsilon = \frac{2}{3}\eta$, the condition for inflation to occur reduces then to $\epsilon \ll 1$.

As we have mentioned above there are two different ways in which inflation may end:

i/- ψ reachs the value ψ_c during the slow-rolling phase, then inflation ends through instability of the field φ and ψ drops to its global minimum.

ii/- If the ϵ grows to be of order unity, the slow-rolling conditions (20) and (21) are no longer valid and inflation ends before ψ reaches ψ_c. This may occur at some value $\psi_e > \psi_c$.

In our model, inflation ends in the second way. Indeed, the value of ψ at which $\epsilon \sim 1$ is by means of Eq.(20)

$$\psi_e \sim O(m_p) \tag{22}$$

which is much greater than ψ_c given by the equation (4), if we take $\mu \sim 10^{15}$ Gev (grand unification scale) and the coupling constant of order 10^{-1} (for particle physics considerations). This is a plausible choice in order to achieve a vacuum-dominated inflation as we shall see below.

After the end of inflation the field ψ undergoes a phase of oscillations about its minimum during which its energy decreases untill the vacuum energy $V = \kappa^2\mu^4$ becomes dominant, then the second stage of inflation begins.

Since we are interested in $\psi < \psi_{60}$ where ψ_{60} is the value of ψ, 60 e-foldings before the end of inflation (see below), we must impose the following condition in order to achieve a vacuum-dominated inflation:

$$\frac{\lambda}{4}\psi_{60}^4 \ll \kappa^2\mu^4 \tag{23}$$

The number of e-foldings of expansion which occur between the two scalar field values ψ_i and ψ_f is given by:

$$N(\psi_i, \psi_f) = -\frac{8\pi}{m_p^2}\int_{\psi_i}^{\psi_f}\frac{V(\psi)}{V'(\psi)}d\psi \tag{24}$$

The number of Hubble times H(k) of which inflation occurs after a given scale leaves the horizon is:

$$N(k) = 62 - \ln\frac{10^{16}\text{ GeV}}{V_I^{\frac{1}{4}}} - \ln\frac{k}{a_0 H_0} \tag{25}$$

where subscript 0 indicates present values [14]. The inflationary energy scale is $V_I^{\frac{1}{4}} \sim 10^{16}$ Gev which makes the Observable Universe leaving the horizon about 60 e-foldings before the end of inflation. The value of ψ_{60} is then calculated by taking $N = 60, \psi_f = \psi_c$ and $\psi_i = \psi_{60}$ in the expression (24):

$$\psi_{60} = \Big[\frac{1}{\kappa^2\mu^2}\Big(\frac{g^2}{2} - \frac{30}{\pi}m_p^2\frac{\lambda}{\mu^2}\Big)\Big]^{-\frac{1}{2}} \tag{26}$$

Substitute (26) in (23) yelds:

$$\lambda \simeq 5\times 10^{-12} \tag{27}$$

3.3 The density perturbation

The adiabatic density perturbation which is generally though to be responsible for large-scale structure originates as a vacuum fluctuation during inflation. Its spectrum is determined by the quantity δ_H which is given in the slow-rolling approximation by the expression [14]:

$$\delta_H^2(k) \ = \ \frac{32}{75}\frac{V_*}{m_p^4}\frac{1}{\epsilon_*} \tag{28}$$

where ϵ is the slow-rolling parameter defined earlier and the subscript $(*)$ indicates that the right-hand side is to be evaluated as the comoving scale k equals the Hubble radius $(k = aH)$ during inflation.

Assuming that gravitationnel waves are negligible, the spectrum measured by COBE is [19]

$$\delta_H \ = \ 1.91\times 10^{-5} \tag{29}$$

Comparing Eq(28) with the value Eq(29) yelds

$$m_p^{-3}\frac{V^{\frac{3}{2}}}{V'} = 5.3 \times 10^{-4} \tag{30}$$

Since during the last stage of inflation we have $V \simeq \kappa^2\mu^4$ [Eq.(23)] this gives:

$$m_p^{-3}\frac{\kappa^3\mu^6}{\lambda\psi^3} \simeq 5,3.10^{-4} \tag{31}$$

The left-hand side of the expression (31) is to be evaluated when the relevent scales cross the horizon ($k = aH$). This occurs about 60 e-foldings before the end of inflation [Eq.(26)]. which corresponds to a non strange value of λ:

$$\lambda \simeq 5.10^{-12} \tag{32}$$

which is very consistent with the value of Eq.(27). In fact, the smallness of λ is a common feature of all models of inflation.

A scalar potential of the form of Eq.(16) has been considered in a nonsupersymmetric context in the reference [20] where the authors have studied the two ways the inflation ends in, and in the simplest case where the scales of astrophysical interest crossed outside the Hubble radius during the second stage of inflation the correct size of density perturbations constraints the parameter λ, to be $\lambda \ll 10^{-12}$.

4 Discussion and conclusion

The above result is really consistent with the condition (27) for the second stage of inflation to take place. Indeed, we can consider that the relevent scales for the structure formation leave the horizon only during the second stage, and so we do not need to impose such a fine-tuning. This consideration agrees with the fact that the inflation that increased the volume of the observable Universe beyond the Hubble radius must occur at an energy scale $V^{\frac{1}{4}} \overset{l}{\sim} 4 \times 10^{16}$ GeV [18], which is true for the vacuum energy that drives the second stage inflation.

From Eq.(22) it is clear that the existence of two stages of inflation does not depend on the parameter λ. This is due to the fact that we have neglected the vacuum energy term, otherwise we would have the condition :

$$\sqrt{\lambda} \geq \frac{8\pi}{3m_p^2}\kappa\mu^2 \tag{33}$$

to have a break of the inflationary phase at some value ψ_+ of the scalar field. But for the inflation to restart we must impose a strong constraint on λ [Eq.(27)]. This required value is confirmed by the constraint imposed by the spectrum of density perturbations [Eq.(32)]. consequently the regime of two-stage inflation is narrower than the simple one.

The idea of two epochs of inflation was also invoked in the reference [17] but in this case the two-stage (double inflation) are drived respectively by two massive non interacting scalar fields :

a heavy one with mass m_h and a light one with mass m_l. The system is described by the scalar potential:

$$V(\phi_l\,,\,\phi_h) \,=\, \frac{1}{2}\Big(m_h\phi_h^2 \,+\, m_l\phi_l^2\Big) \tag{34}$$

A double inflation with a break (The two phases of inflation are separated by a non inflationary one with $a(t) \sim t^{2/3}$) occurs when the energy of the light field is too small to drive inflation at the moment when $\phi_l \,=\, m_p/2\pi$, and the energy of the heavy one cannot sustain inflation any more. This gives the condition :

$$m_l\phi_l \ll mhm_p/2\pi \tag{35}$$

this model contains two free parameters m_l and m_h which are fundamental.

In this paper we have considered an extension of the usual superpotential of the supersymmetric hybrid inflation. This model describes two disconnected stages of inflation.The first is a simple chaotic inflation while the second one is a vacuum-dominated one. The typical scale of the first stage is of order of the Planck mass. After the end of this stage, the inflaton moves towards the minimum of the potential around which it begins to oscillate. During this intermediate stage the energy density is reduced through expansion. The second stage begins when the energy density falls below the false-vacuum energy density associated with a phase transition involving the other field. The homogeneity beyond the Hubble radius that was produced during the first stage makes natural the onset of the second stage.

In this model a sever fine-tuning of the self-coupling constant is required as a condition for the second stage of inflation to occur. But the same fine-tuning is also required in order to allow the correct magnitude of the spectrum of density perturbations. Therfore, even though the fine-tuning is always seen as an unplausible feature of inflationary models, this fact seems to give a self consistency for this model. Furthermore, the fine-tuning usually imposed by the COBE observations on the self-coupling constant of the chaotic inflation model is somewhat relaxed in this model. Indeed, it is not necessary to require the generation of the correct spectrum of densty fluctuations during the first stage since it will be sufficient to assume that the relevant scales for the structure formation leave the horizon only during the second stage of inflation. The role of the first stage in this model can then be restricted to provide the necessary homogeneity for the onset of the second one.

The fact that the two stages occur at two different energy scales, namely the Planck and the GUT scales can be seen as one of the advantages of this extension. This exhibits the succession of the phase transitions in the early Universe.

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