Natural Inflation in SUSY and Gauge-Mediated Curvature of the Flat Directions

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Abstract
Supersymmetric theories often include the non-compact directions in the field space along which the tree level potential grows only up to a certain limited value (determined by the mass scale of the theory) and then stays constant for the arbitrarily large expectation value of the field parametrizing the direction. Above the critical value, the tree-level curvature is large and positive in the other directions. Such plateaux are natural candidates for the hybrid inflaton. The non-zero $F$-term density along the plateau spontaneously breaks SUSY and induces the one-loop logarithmic slope for the inflaton potential. The coupling of the inflaton to the Higgs fields in the complex representations of the gauge group, may result in a radiatively induced Fayet–Iliopoulos $D$-term during inflation, which destabilizes some of the squark and slepton flat directions. Corresponding soft masses can be larger than the Hubble parameter and thus, play a crucial role for the Affleck–Dine baryogenesis.

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Introduction

The usual problem with inflationary scenarios is that they require introduction of the unmotivated small parameters: the scalar field with an extremely small self-coupling and the flat potential. In the non-supersymmetric theories the scalars with such a flat potential do not emerge naturally. On the contrary, it is well known that the supersymmetric theories often admit the flat vacuum directions, which are flat to all orders in perturbation theory in the unbroken SUSY limit. They gain a very small curvature (typically $\sim 100$ GeV), induced by supersymmetry breaking and, therefore, inflation along these directions would be marked with a very small value of the Hubble constant. Such an inflation can probably, at best, be considered on top of the earlier inflation with a sufficiently large Hubble parameter ($\sim 10^{13}$ GeV or so) motivated by the large density perturbations consistent with COBE data. However, even if not responsible for the inflation, the supersymmetric flat directions can still play a very important role in cosmology.

The aim of the present paper is twofold: first, to point out some natural candidates for the inflation that involve a large mass scale (perhaps motivated from the particle physics), e.g. such as the GUT scale $M_G$, and secondly to discuss the behaviour of the flat directions during such an inflation. We will argue that supersymmetric theories often include certain ‘plateau’ directions in the field space, parametrized by the scalar fields $X$, which in the global

1Unless the inflaton is a pseudo-Goldstone particle.

2Here and below we are talking about the vacuum energy driven inflation. The situation can be very different in the case of the kinetic inflation, e.g. in string cosmology.
minimum have large positive curvature induced by large VEVs of the other field(s) \( \phi \). Nevertheless, the potential along such a direction grows up to a certain value and then becomes frozen, since for large values of \( X \) the \( \phi \) VEV is switched-off and the tree-level potential is exactly flat for \( X \to \infty \). Since for a large \( X \) the potential is non-zero, supersymmetry is broken and there is a one-loop induced effective potential that drives \( X \) back to the SUSY minimum. This scenario leads to a natural realization of the ‘hybrid inflation’ idea invented by Linde [6].

During inflation the flat directions, which have no tree-level couplings with an inflaton, gain masses \( \sim H \) through the gravity-mediated supersymmetry breaking [7]. The gauge-charged flat directions, however, can gain the gauge-mediated soft masses if the inflaton interacts with some gauge-non-singlet fields [8]. It is well known [10] that the gauge interaction also can be a messenger of SUSY breaking, and in fact more efficient than gravity, provided the messenger scale is lower than \( M_{\text{Planck}} \). This is precisely what is going to happen in the inflationary scenario discussed above, since the inflaton has renormalizable couplings with gauge nonsinglet fields \( \phi \) and itself also can be a gauge non-singlet field, e.g. can carry colour and electric charge. Thus, during inflation the \( \phi \)-fields play the role of messengers of supersymmetry breaking and radiatively induce the universal (up to charges) two-loop soft masses for all the gauge non-singlet flat directions, which have no tree-level couplings with the inflaton. The important thing is that these soft masses can be greater than \( H \), typical magnitude of the gravitationally induced soft masses. Here we observe that when the inflaton couples to the fields in the complex representations, there is a one-loop-induced Fayet–Iliopoulos
D-term during inflation. This term can dominate both the gauge-mediated two-loop and the gravity-mediated soft terms, and destabilize some of the squark/slepton flat directions during inflation. Therefore it can play a crucial role for the Affleck–Dine mechanism of baryogenesis [11].

Simple Examples

The simplest inflationary scenario with the above properties was considered in [12]. The superpotential of the model is

\[ W = \frac{1}{2} f X \phi^2 - X \mu^2, \]  

(1)

where \( \phi \) is a superfield that breaks the GUT symmetry and \( \mu \) is a mass scale such that \( \mu \sqrt{f} = M_G \); \( \phi \) can be a component of the Higgs field in the real (adjoint) or complex representation, e.g. a spinor of \( SO(10) \) or a 6-plet of \( SU(6) \) (see below). The above superpotential is the most general one compatible with the GUT symmetry and an \( R \)-symmetry under which \( X \) carries one unit of the superpotential charge. The scalar potential is given by

\[ V = \left| \frac{1}{2} f \phi^2 - \mu^2 \right|^2 + f^2 |X|^2 |\phi|^2 + D - \text{terms} \]  

(2)

This theory has a unique supersymmetric vacuum with \( \phi^2 = \frac{2}{f} \mu^2 \) and \( X = 0 \). However, minimization with respect to \( \phi \) for the fixed values of \( X \) shows that for \( X > X_c = \frac{\mu}{\sqrt{f}} \), the minimum is at \( \phi = 0 \), the potential is flat in the \( X \) direction and has a large curvature (\( \sim f |X| \)) in the \( \phi \) direction.

3A similar superpotential was considered in [13], but the inflaton one-loop effective potential, which plays a crucial role, was ignored there.
This system, under the assumption of the chaotic initial conditions with $|X| \gg X_C$ naturally leads to the inflation.

Since the curvature in the $\phi$ direction is very large, we expect that $\phi$ will rapidly settle in its instant minimum with $\phi = 0$. In contrast, the curvature in the $X$ direction is zero and the system will evolve towards the global minimum very slowly. This state is dominated by large $|F_x| = \mu^2$ term, which leads to the inflation. The one loop-corrections

$$\Delta V = \left(\frac{-1}{64\pi^2}\right)F \tau M^4 \ln \frac{M^2}{\Lambda^2}$$

(3)

provide non-zero curvature driving $X$ towards the SUSY vacuum. The one-loop corrected effective potential for the large $X$ behaves as

$$V = \mu^4 \left(1 + \frac{f^2}{16\pi^2} \left[2\ln \frac{f^2|X|^2}{\Lambda^2} + 3\right]\right).$$

(4)

The phase transition with gauge symmetry breaking takes place only after the $X$ field drops to its critical value $X_c$. Below this point, all the VEVs rapidly adjust to their supersymmetric values.

Since the $F_x$-term, which dominates the inflationary Universe, splits the Fermi–Bose masses of the gauge-non-singlet superfield(s) (in the above case $\phi$), there are radiatively induced two-loop soft masses of all gauge-non-singlet flat directions (and in particular squarks and sleptons). In this case $\phi$ plays the role of the messenger of the inflaton-induced SUSY breaking for all other gauge-non-singlet fields. Integrating out the heavy messengers at each point of inflationary trajectory, we end up with a two-loop ($X$-dependent) soft masses for the light scalars

$$m^2 \sim \left(\frac{\alpha}{4\pi}\right)^2 \frac{\mu^4}{|X|^2}.$$ 

(5)
These masses can be larger than the Hubble constant:

\[ H^2 = \frac{\mu^4}{3M^2} \]  

(6)

where \( M = \frac{M_{\text{Planck}}}{\sqrt{8\pi}} \). Thus, gauge-mediated corrections can be the dominant source of the scalar soft masses during inflation and cannot be neglected.

The sign of the soft masses plays a crucial role for the baryogenesis via the Affleck-Dine mechanism, since this mechanism requires large expectation values of squarks and sleptons along the flat directions.

Here we wish to show that in more generic (and realistic) cases there are similar one-loop corrections with either sign. They appear due to the one-loop-induced Fayet–Iliopoulos \( D \)-term during inflation. This generically happens in the theories in which the inflaton couples to the Higgs fields in the complex representation that is required for lowering the rank of the group in all GUTs other than \( SU(5) \). These corrections tend to destabilize the flat directions during inflation and, thus, naturally lead to the Affleck–Dine mechanism of the baryogenesis.

Let us consider a minimal structure that leads to such a picture. We introduce a pair of Higgs fields \( \phi^+, \phi^- \) with opposite charges under a certain \( U(1) \)-group. We will think of this \( U(1) \) as being a broken Abelian subgroup of some GUT symmetry under which \( \phi^+, \phi^- \) transform in the complex representations. We assume that in the tree-level global minimum \( U(1) \) is broken by the VEV of \( \phi^+, \phi^- \) fields (along the \( D \)-flat direction \( \phi^+ = \phi^- = M_G \)) triggered by the singlet \( S \). Furthermore, we assume that this VEVs give mass to some charged states \( A^-, A^+ \) by mixing them with the neutral one \( X \). The
simplest superpotential with above properties can be chosen as

$$W = f S \phi_+ \phi_- + \mu^2 S + X(a_+ A_+ \phi_- + a_- A_- \phi_+).$$

(7)

On top of this we assume that there are some U(1)-charged flat directions that do not have direct couplings with the above fields in the superpotential. Again, this superpotential is the most general under $U(1)$ and the $R$-symmetry under which $S$ carries one unit of the superpotential charge.

The scalar potential of this system reads

$$V = \left| f \phi_+ \phi_- - \mu^2 \right|^2 + |a_+ A_+ \phi_- + a_- A_- \phi_+|^2 + |a_+ X \phi_-|^2 + |a_- X \phi_+|^2$$

$$+ \left| f S \phi_+ + a_+ A_+ X \right|^2 + \left| f S \phi_- + a_- A_- X \right|^2$$

$$+ \frac{g^2}{2} \left( |\phi_+|^2 - |\phi_-|^2 + |A_+|^2 - |A_-|^2 + q_i |Q_i|^2 \right)^2,$$  

(8)

where the $Q_i$ are all other charged fields and $q_i$ are their charges. This potential has a supersymmetric minimum with

$$\phi_+ = \phi_- = \frac{\mu}{\sqrt{f}} \quad S = X = A_+ = A_- = 0.$$  

(9)

In addition, we assume that there are flat directions along which some combinations of the $Q_i$ fields are undetermined. In this vacuum one combination $\frac{a_+ A_+ + a_- A_-}{\sqrt{a_+^2 + a_-^2}}$ is mixed with $X$ and gets mass $\sqrt{\frac{a_+^2 + a_-^2}{f}} \mu$. Minimizing the potential for the different fixed values of $X$ we find that for $|X| > X_c = \sqrt{\frac{a_+}{a_-}} \mu$ the minimum is at $A_+ = A_- = \phi_- = \phi_+ = 0$ and $S$ undetermined. Thus, the potential has a fixed value $\mu^2$ and a vanishing curvature in the $X$ direction for any $|X| > X_c$. In addition the curvature in the $S$ direction is also zero, since $F_s$ breaks SUSY and the scalar component is a partner of the Goldstone fermion. As before, the one-loop corrections provide logarithmic slope for the
\( \text{and } S \text{ and drive them towards the global minimum. For simplicity we will quote the form of the effective potential along the } S = 0. \text{ For } S = 0 \text{ and } |X| \gg X_c \text{ the contribution to the one-loop effective potential only comes from the superfields } \phi_-, \phi_+ \text{ whose masses are split by the non-zero } F_s \text{-term. This superfield delivers two massive fermions with masses } a_+|X| \text{ and } a_-|X| \text{ and two complex scalars with mass-squared }

\[
m^2_\pm = \frac{|X|^2}{2} \left[ (a_+^2 + a_-^2) \pm (a_+^2 - a_-^2) \sqrt{1 + \frac{4f^2\mu^4}{|X|^4(a_+^2 - a_-^2)}} \right]. \tag{10}
\]

For the large \( X \) the one-loop-corrected effective potential behaves as

\[
V = \mu^4 \left( 1 + \frac{f^2}{16\pi^2} \left[ 1 + 2 \left( \frac{a_+^2 \ln a_+^2 - a_-^2 \ln a_-^2}{(a_+^2 - a_-^2)^2} + 2\ln \frac{|X|^2}{\Lambda^2} \right) \right] \right). \tag{11}
\]

The crucial difference between this scenario and the one discussed before is that now, during inflation, one-loop radiative corrections induce the Fayet–Iliopoulos \( D \)-term for the \( U(1) \) gauge superfield, which will be compensated by some of the flat directions in the \( Q_i \) sector that carry the \( U(1) \) charge of the appropriate sign. The reason for the appearance of this \( D \)-term is that during inflation the messengers \( \phi_+ \) and \( \phi_- \) have different masses, since \( a_+^2 \neq a_-^2 \), and their contributions do not cancel out in the one-loop diagram (see Fig.1). In the leading order, the corresponding soft masses for the light scalars are proportional to

\[
q_i \left( \frac{\alpha}{4\pi} \right) \frac{f^2\mu^4}{|X|^2} \left[ \frac{a_+^2 + a_-^2}{(a_+^2 - a_-^2)^2} \ln \frac{a_+^2}{a_-^2} + \frac{2}{a_+^2 - a_-^2} \right] \sim H^2 \frac{M^2}{|X|^2} \left( \frac{\alpha}{4\pi} \right). \tag{12}
\]

Now, the inflation ends when the slow roll conditions break down, that is when either \( \left| \frac{M V'}{V} \right| \sim 1 \) or \( |V''| \sim H^2 \), where the prime denotes the derivative in the inflaton direction. For the inflaton potential given by (11), this
One-loop contribution to the curvature of the flat direction $Q$.

happens when $X$ drops to the value $|X| \sim \sqrt{CM}$, where $C \sim \frac{f^2}{8\pi^2}$ is a loop factor. Thus, at the end of the inflation the gauge-mediated soft masses are $\sim H^2C^{-1}\left(\frac{\alpha}{4\pi}\right)$. In the interval $X_c < |X| < \sqrt{CM}$, between the end of inflation and the gauge-symmetry-breaking phase transition, the ratio of the gauge-mediated to gravity-mediated soft masses is

\[
\frac{m_{gauge}^2}{m_{gravity}^2} \sim \frac{M^2}{|X|^2}\left(\frac{\alpha}{4\pi}\right). \tag{13}
\]

We see that (at least on the later stages of the inflation) the gauge-mediated corrections are dominant and can thus play the crucial role for the Affleck-Dine scenario of the baryogenesis.
Realistic Example

The inflationary scenario outlined above can naturally emerge in many realistic GUTs in which the Higgs fields in the complex representation are lowering the rank of the group. Here we consider one example motivated from particle physics. This is an $SU(6)$ GUT in which the doublet–triplet splitting problem is solved by the pseudo-Goldstone mechanism [15]. The minimal Higgs structure, group theoretically required for the breaking of $SU(6)$ down to $G_W = SU(3)_c \otimes SU(2)_L \times U(1)_Y$ is one adjoint 35-plet $\Sigma_i^k$ and a pair of fundamental and anti-fundamental representations $(6, \bar{6}$-plets) $\phi_i$ and $\bar{\phi}^i$ ($i, k = 1, 2, \ldots, 6$ are $SU(6)$ indexes). This minimal structure suffices to deliver a pair of light electroweak Higgs doublets and naturally solve the doublet–triplet splitting problem. The crucial assumption is that $\Sigma$ and $\phi, \bar{\phi}$ have no direct cross couplings, in the superpotential, so that the Higgs superpotential has the form:

$$W_{Higgs} = W_\Sigma(\Sigma) + W_\phi(\phi, \bar{\phi}).$$

(14)

The absence of the possible cross term $\bar{\phi}\Sigma\phi$ can be guaranteed by the exact discrete or continuous symmetries of the theory, which we will not specify here; just note that the one possible candidate is an $R$-symmetry (for the concrete realizations see [15],[16]). For the breaking of the gauge group to $G_W$ the $\Sigma$ and $\Phi$-fields should pick up the VEVs in the $SU(4) \otimes SU(2) \otimes U(1)_q$-invariant and $SU(5)$-invariant directions respectively:

$$\Sigma = \text{diag}(1, 1, 1, -2, -2) \sigma \quad \phi = (\phi_+, 0, 0, 0, 0, 0) \quad \bar{\phi} = (\phi_-, 0, 0, 0, 0, 0).$$

(15)
Due to the absence of the cross couplings, the Higgs superpotential has an $SU(6)_{\Sigma} \otimes SU(6)_{\phi}$ global symmetry, which gets broken to

$$G_{\text{global}} = [SU(4) \otimes SU(2) \otimes U(1)_{q}]_{\Sigma} \otimes [SU(5)]_{\phi}. \quad (16)$$

Simple counting of the Nambu–Goldstone modes shows that there is a pair of massless (in the SUSY limit) electroweak doublet states, which are not eaten up by the gauge fields and are physical Higgs particles.

In this model quarks and leptons are placed in the $15^\alpha + \bar{6}^\alpha + \bar{6}'^\alpha$ representations ($\alpha = 1, 2, 3$ is a family index), which under the $SU(5)$ group decompose as

$$10^\alpha + 5^\alpha + \bar{5}^\alpha + 5'^\alpha + 1'^\alpha + 1^\alpha \quad (17)$$

and we see that there are on top of the usual chiral matter (quarks and leptons in three $10 + \bar{5}$ copies), extra vector-like states $5 + \bar{5}$ and two singlets per family (one combination of singlets can play the role of the right-handed neutrino). These extra states gain GUT scale masses from the VEV of $\phi, \bar{\phi}$.

The fermion masses in this GUT were studied in detail in [17] and the interested reader is referred to this paper. Here we will concentrate on the part of the superpotential that is responsible for the breaking $SU(6) \rightarrow SU(5)$ and for the generation of the heavy matter masses. The simplest possible form for this sector can be chosen as

$$W = f S \phi \bar{\phi} - S \mu^2 + a_{\alpha} \phi 15^\alpha \bar{6}^\alpha + b_{\alpha} \phi \bar{6}'^\alpha X_{\alpha}. \quad (18)$$

where $X_{\alpha}$ are three additional singlets. Here we have introduced a matter parity under which only matter multiplets change sign and another $Z_2$ symmetry under which only $\bar{6}'$ and $X$ superfields change sign. Without loss
of generality, we have diagonalized Yukawa coupling matrices by the field redefinition. The potential of this system reads

\[ V = \left| f\phi\bar{\phi} - \mu^2 \right|^2 + a_\alpha^2 \left| \bar{6}^\alpha - \bar{\phi}^\alpha \phi \right|^2 + a_\alpha^2 \left| \phi^2 X_\alpha \right|^2 + |\phi\bar{6}^\alpha|^2 \]

\[ + \left| fS\bar{\phi} + b_\alpha \bar{6}^{\prime\alpha} X_\alpha \right|^2 + |fS\phi + a_\alpha 15^\alpha \bar{6}^\alpha|^2 + D - \text{terms} \quad (19) \]

(SU(6) indices are suppressed). Let us investigate the behaviour of the above potential along the non-compact $D$-flat direction

\[ 15^1_{ik} = X(\delta^2_i \delta^2_k - \delta^2_i \delta^2_k) \quad \bar{6}^{1i} = X\delta^2_i \quad \bar{6}^{2i} = X\delta^1_i \quad X_2 = X \quad (20) \]

where $X$ is a parameter. It is not difficult to notice that this direction is not $F$-flat and has a large $\sim \mu^2$ positive curvature at the origin $X = 0$. However, for

\[ |X|^2 > X_c = \max \left( \frac{f\mu^2}{a_1 b_2}, \frac{f\mu^2}{\sqrt{2a_2 b_2}} \right) \quad (21) \]

the curvature vanishes for arbitrarily large values of $X$, and the tree-level potential stays constant $V_o = \mu^4$. This is because the instant minimum in all other fields (coupled to the given direction) is at their zero value as far as $|X| > X_c$. The vanishing of the $\phi, \bar{\phi}$ fields, which gain large $\sim |X|^2$ positive masses, leads to the non-zero cosmological constant that drives inflation. During this period supersymmetry is spontaneously broken by the expectation value $F_s = \mu^2$, which splits masses of the $\phi, \bar{\phi}$ superfields and induces the one-loop effective potential similar to the one of (11). We see that the inflationary scenario discussed in the previous section emerges naturally in this context. Note that $D$-flatness conditions do not require $X_2 = |X|$, so that
we can treat $X_2$ as a second parameter of the plateau. The critical values above which the curvature vanishes will then be defined by the condition

$$|XX_2| > \max \left( \frac{f\mu^2}{a_1b_2}, \frac{f\mu^2}{\sqrt{2a_2b_2}} \right)$$

(22)

and the inflationary trajectory will be parametrized by both $X$ and $X_2$. The $D$-flat direction we have chosen is in no way unique and one may equally well choose in (20) other combinations of $15^\alpha, \bar{6}^\beta, \bar{6}^\gamma$ and $X_\nu$ with $\alpha \neq \beta$ and $\gamma \neq \nu$. The simple rule is that, from each operator, only one VEV can participate in the plateau direction, since in the opposite case some of the $F$-terms can grow unbounded along this direction.

Now what about the gauge-mediated soft masses during inflation? To answer this question, first note that an interesting difference between this inflation and the simple toy model discussed in the previous section is that the inflaton parametrizes the gauge-nonsinglet direction along which the colour and the electric charge are broken. *Per se*, this offers an interesting scenario of baryogenesis through the out-of-equilibrium decay of the inflaton. To see which is the unbroken gauge group during inflation we have to know the relative orientation of the $\Sigma$ VEV in the basis in which other VEVs are given by (20). Precise orientation will be decided by the radiative corrections, which induce an effective potential for the pseudo-Goldstone modes. The detailed analysis of this dynamics will not be attempted here and, for definiteness, we assume that the form of sigma is the one given by (15). In such a case the unbroken gauge symmetry in the inflationary epoch would be

$$G_{inf} = SU(2)_C \otimes SU(2)_L \otimes U(1)_{YC}$$

(23)
where $SU(2)_C$ is a colour subgroup and the generator of $U(1)_Y$, $\text{diag}(0, 0, 1, 1, -1, -1)$, is a combination of hypercharge and colour. It is expectable that some of the squark and slepton fields which are not coupled, at tree level, with the inflaton direction, will be destabilized by one-loop-induced soft masses proportional to their charges under the broken generators. Note that, in this particular example, there is no contribution to one-loop soft masses from the $U(1)_Y$ gauge fields. This is because in the above case $\phi$ and $\bar{\phi}$, which play the role of the SUSY-breaking messengers during inflation, are decoupled at tree level from the $\Sigma$, and the masses of their $U(1)_Y$-charged components are degenerate. In contrast, the masses of the components that carry charge under broken generators (e.g. $\text{diag}(-5, 1, 1, 1, 1, 1)$) are not degenerate. This can be seen from the tree-level values of these mass terms along the inflationary trajectory:

$$
|\bar{\phi}_1|^2a_1^2|X|^2 + |\bar{\phi}_2|^2(a_1^2 + 2a_2^2)|X|^2 + |\phi_k|^22a_2^2|X|^2 + |\phi_1|^2b_2^2|X_2|^2 + |\phi_2|^2(b_1^2|X|^2 + b_2^2|X_2|^2) + |\phi_k|^2b_2^2|X_2|^2 + f\bar{\phi}\phi\mu^2 + \text{h.c.}
$$

(24)

where $k > 2$. The resulting one-loop contribution to the soft masses in general is $\sim \left(\frac{\alpha}{\pi}\right) \frac{\mu^4}{|X|^2}$.

**Conclusions**

The aim of the present letter was to point out that the supersymmetric theories that include a large mass scale may contain a potential source for the natural inflation—non-compact plateau directions in the field space along which the tree-level potential has a vanishing curvature and large constant value. Such directions may be parametrized by the gauge-singlet or gauge-
non-singlet fields. During inflation at least some of the chiral superfields that have tree-level couplings with the inflaton suffer from the Fermi–Bose mass splitting and effectively play the role of messengers for the gauge-mediated supersymmetry breaking. In the case of the complex representations, this breaking can generate the one-loop Fayet–Iliopoulos $D$-term during inflation, which can destabilize some of the squark/slepton VEVs. The resulting gauge-mediated soft masses are typically larger than the gravity-mediated ones ($\sim H$); they can thus play a crucial role for the Affleck–Dine scenario of baryogenesis.

Acknowledgements

I learned that recently A.Linde and A.Riotto have discussed in a different context the relevance of the SUSY breaking at the end of inflationary epoch [18].

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