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## Minimalism in Inflation Model Building <sup>\*</sup>

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### Abstract

In this paper we demand that a successful inflationary scenario should follow from a model entirely motivated by particle physics considerations. We show that such a connection is indeed possible within the framework of concrete supersymmetric Grand Unified Theories where the doublet-triplet splitting problem is naturally solved. The Fayet-Iliopoulos  $D$ -term of a gauge  $U(1)_\xi$  symmetry, which plays a crucial role in the solution of the doublet-triplet splitting problem, simultaneously provides a built-in inflationary slope protected from dangerous supergravity corrections.

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## I. INTRODUCTION

The flatness and the horizon problems of the standard big bang cosmology are elegantly solved if during the evolution of the early Universe the energy density happens to be dominated by some vacuum energy and comoving scales grow quasi-exponentially [1]. An inflationary stage is also required to dilute any undesirable topological defects left as remnants after some phase transition taking place at early epochs. The vacuum energy driving inflation is generally assumed to be associated to some scalar field, the inflaton, which is displaced from the minimum of its potential. As a by-product, quantum fluctuations of the inflaton field may be the seeds for the generation of structure formation.

There are many problems one has to face in building up a successful inflationary model. First of all, the level of density and temperature fluctuations observed in the present Universe,  $\delta\rho/\rho \sim 10^{-5}$ , require the inflaton potential to be extremely flat. This is in contrast with the requirement that the couplings of the inflaton field to other degrees of freedom cannot be too small otherwise the reheating process, which converts the vacuum energy into radiation at the end of inflation, takes place too slowly: large couplings induce large loop corrections to the inflaton potential, spoiling its flatness. Introducing very small parameters to ensure the extreme flatness of the inflaton potential seems very unnatural and fine-tuned in most non-supersymmetric theories. However, this technical naturalness may be achieved in supersymmetric models [2] because the nonrenormalization theorem guarantees that the superpotential is not renormalized to all orders of perturbation theory [3]. The perturbative renormalization of the Kähler potential, however, can be crucial for the inflationary dynamics due to a non-zero energy density which breaks supersymmetry spontaneously during inflation [4,5]<sup>1</sup>, independently whether this energy density is an input or results from some strong dynamics [6,5]. Secondly, there is the (aesthetic) problem of embedding a successful inflationary scenario in the framework

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<sup>1</sup>In particular, this renormalization can be due to some strongly coupled particles whose condensate generates the inflaton scale dynamically [5].

of some well-motivated particle physics models.

To our opinion, one should apply a sort of "minimal principle" [7] requiring that any successful inflationary scenario should naturally arise from models which are entirely motivated by particle physics considerations and should not involve (usually complicated and *ad hoc*) sectors on top of the existing structures. Recently such attempts have been made in [5], in the framework of dynamical grand unified symmetry breaking, and in [7] where the inflaton candidates were identified in some models of gauge-mediated supersymmetry breaking.

It is the main purpose of this paper to demonstrate a possibility of the connection between the inflationary scenario and the particle physics problems, within the framework of concrete Grand Unified Theories. In doing that, we will be entirely motivated by the solution to a serious problem arising in supersymmetric Grand Unified Theories (SUSY GUTs), namely the doublet-triplet splitting problem. We will show that the model which is able to solve this problem also naturally incorporates a built-in inflationary scenario. We will also show that our proposal escapes the usual slow-roll problems posed by supergravity corrections in  $F$ -term dominated inflation. The supergravity corrections usually induce large (of order the Hubble parameter  $H$ ) curvature for the inflaton slope and inflation does not take place [8]. The appearance of such a large curvature reflects the fact that SUSY must be broken during inflation. This mass does not disappear in the limit in which the Planck mass  $M_{Pl}$  tends to infinity when  $H$  is held fixed. As was suggested in [9], one possible way out to avoid this problem is to have inflation dominated by a  $D$ -term <sup>2</sup>. Indeed, in the de Sitter space the gravity-transmitted  $D$ -type supersymmetry breaking can be much weaker than the  $F$ -type counterpart and the slow-roll problem may be avoided. Large  $D$ -term driving inflation can be induced, for example, if the theory contains a gauge  $U(1)_\xi$  factor with a nonvanishing Fayet-Iliopoulos  $D$ -term

$$\int d^4\theta \xi V. \tag{1}$$

This term may be present in the underlying theory from the very beginning ( it is allowed

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<sup>2</sup>See also comment in [10].

by a gauge symmetry, unless  $U(1)$  is embedded in some non-Abelian group<sup>3</sup>) or may appear in the effective theory after some heavy degrees of freedom have been integrated out. Moreover, it looks particularly intriguing that an anomalous  $U(1)_\xi$  symmetry is usually present in string theories [12]<sup>4</sup> and the anomaly cancelation is due to the Green-Schwarz mechanism [13]. The corresponding Fayet-Iliopoulos term is given by

$$\xi = \frac{g^2}{192\pi^2} \text{Tr}\mathbf{Q} M^2, \quad (2)$$

where  $M = M_{Pl}/\sqrt{8\pi}$  is the reduced Planck mass and  $\text{Tr}\mathbf{Q} \neq 0$  indicates the trace over the  $U(1)_\xi$  charges of the fields present in the spectrum of the theory.

On the other hand, the anomalous  $U(1)_\xi$  can play also a crucial role in the solution of the doublet-triplet splitting problem [15]. It is therefore natural to attempt to reconcile these two implications coming from theories containing an anomalous  $U(1)_\xi$  symmetry and to construct a model that would solve the hierarchy problem and simultaneously predict a successful stage of inflation in the early universe. Before proceeding, we would like to point out that in our scenario the use of an anomalous  $U(1)_\xi$  is not strictly necessary. What is really crucial is the presence of a gauge  $U(1)_\xi$  with nonvanishing  $D$ -term (1). In this respect any gauge  $U(1)_\xi$  would be suitable for our purposes, but the advantage of an anomalous  $U(1)_\xi$  is that  $\xi$  is not an input parameter but is fixed from the expression (2). Therefore, we keep our discussion as general as possible and explicitly indicate the difference between an anomalous and nonanomalous  $U(1)_\xi$  when the difference is important.

Our inflationary scenario can be regarded as a realistic variant of hybrid inflation [16]. Typically in this scenario the inflaton field is represented by a gauge singlet coupled to the Higgs field that triggers the end of inflation via a *non-thermal* phase transition with symmetry breaking. Dangerous topological defects (*e.g.* magnetic monopoles in the

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<sup>3</sup> $\xi = 0$  can be enforced by charge conjugation symmetry [11] which flips all  $U(1)$  charges. Such symmetry is possible in nonchiral theories.

<sup>4</sup>Some cosmological implications of the anomalous  $U(1)$  were studied in a different context [14].

grand unified context) may be produced. What is unusual in our scenario is that inflaton is not a gauge singlet, but resides in the component of the adjoint Higgs that breaks GUT symmetry. Consequently the GUT symmetry is broken both during and after inflation, and no monopoles are produced.

## II. THE MODEL

### A. Higgs Sector and the Doublet-Triplet Splitting

Let us briefly describe the main features of the model we have in mind to solve the doublet-triplet splitting problem. It is essentially based on the mechanism of [17]. The novelty in our case is that we incorporate  $D$ -term in the spirit of [17] in order to generate VEVs and therefore simplify the structure of the superpotential.

Let us consider an  $SU(6)$  supersymmetric GUT with one adjoint Higgs  $\Sigma$  and a number of fundamental Higgses  $H_A, \bar{H}^A, H'_A, \bar{H}'^{A'}$ . We assume that each of these fundamentals transforms as a doublet of a certain custodial  $SU(2)_c$  symmetry that is required to solve the hierarchy problem [17]. The index  $A = 1, 2$  is the  $SU(2)_c$ -index. We also assume that  $H_A, \bar{H}^A$  carry unit charges opposite to  $\xi$  and are the ones that compensate  $U(1)_\xi$   $D$ -term in the present Universe. In the context of string inspired anomalous  $U(1)_\xi$  this would simply mean that they carry charges opposite to total trace  $\text{Tr}\mathbf{Q}$ .

The superpotential reads

$$W = c\text{Tr}\Sigma^3 + (\alpha\Sigma + aX + M)H_A\bar{H}'^A + (\alpha'\Sigma + a'X + M')H'_A\bar{H}^A. \quad (3)$$

Minimizing both the  $D$ - and the  $F$ -terms we get the following supersymmetric vacuum which leaves  $SU(3)_c \otimes SU(2)_L \otimes U(1)$  as unbroken gauge symmetry

$$H_{Ai} = \bar{H}^{Ai} = \delta_{A1}\delta_{i1}\sqrt{\frac{\xi}{2}}, \quad H'_A = \bar{H}'^{A'} = 0, \\ \Sigma = \frac{aM' - a'M}{a'\alpha - \alpha'a}\text{diag}(1, 1, 1, -1, -1, -1), \quad X = -\frac{\alpha M' - \alpha' M}{a'\alpha - \alpha'a}. \quad (4)$$

Here  $i, k = 1, 2, \dots, 6$  are  $SU(6)$  indexes. The role of the  $\Sigma$  VEV is crucial since it leaves the unbroken  $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$  symmetry, consequently it can cancel masses of

all upper three or lower three components of the fundamentals [18]. The fundamental VEVs are  $SU(5)$  symmetric, so that the intersection gives the unbroken standard model symmetry group.

In this vacuum the electroweak Higgs doublets from  $H_2, \bar{H}^2, H'_2, \bar{H}'^2$  are massless. This is an effect of custodial  $SU(2)_c$  symmetry. Indeed, since  $H_1$  and  $\bar{H}^1$  break one of the  $SU(3)$  subgroups to  $SU(2)_L$ , their electroweak doublet components become eaten up Goldstone multiplets and cannot get masses from the superpotential due to the Goldstone theorem. This forces the VEVs of  $\Sigma$  and  $X$  to exactly cancel their mass terms and those of  $H_2, \bar{H}^2, H'_2, \bar{H}'^2$  due to the custodial symmetry. This solves the doublet-triplet splitting problem in a natural way.

An alternative possibility would be to relax the requirement of  $SU(2)_c$  custodial symmetry and instead to introduce a number of singlets  $X_A, X'_A$  coupled to the different pairs, as it was suggested by Barr [19]. In this case one has to assume a nonzero VEV for all  $H_A, \bar{H}^A$  fields. Then the doublet masses will be cancelled by singlets just as in our example.

## B. Fermion Masses

Quarks and leptons of each generation are placed in a minimal anomaly free set of  $SU(6)$  group: 15-plet plus two  $\bar{6}_A$ -plets per family. We assume that  $\bar{6}_A$  form a doublet under  $SU(2)_c$  so that  $A = 1, 2$  is identified as  $SU(2)_c$  index <sup>5</sup>. The fermion masses are then generated as in ref. [17] through the couplings ( $SU(6)$  and family indices are suppressed)

$$\bar{H}^A \cdot 15 \cdot \bar{6}_A + \epsilon^{AB} \frac{H_A \cdot H_B}{M_\xi} 15 \cdot 15, \quad (5)$$

where  $M_\xi$  has to be understood as the mass of order  $\sqrt{\xi}$  of integrated-out heavy states (the simplest possibility is to use the 20-plet transforming as doublet under custodial

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<sup>5</sup>Note that  $15 + \bar{6}_A$  just form a fundamental 27-plet of  $E_6$  if  $SU(6) \otimes SU(2)_c$  is viewed as one of its maximal subgroups.



$SU(2)_c$ ). In the case of anomalous  $U(1)_\xi$  the relative charges of the matter field must be fixed from the Green-Schwarz anomaly cancelation. When the large VEVs of  $H_1$  and  $\bar{H}^1$  are inserted, the additional, vectorlike under  $SU(5)$ -subgroup, states:  $\bar{5}$ -s from  $15$ -s and  $\bar{5}$ -s from  $\bar{6}_1$ , become heavy and decouple. Low energy couplings are just the usual  $SU(5)$ -invariant Yukawa interactions of the light doublets from  $H_2$  and  $\bar{H}^2$  with the usual quarks and leptons.

### III. $D$ -TERM DRIVEN INFLATION

Let us now show that model briefly described in the previous section has a built-in inflationary trajectory in the field space along which all  $F$ -terms are vanishing and only the associated  $U(1)_\xi$   $D$ -term is nonzero. As said in the introduction, this peculiar feature will allow inflation to take place without suffering from the slow-roll problem induced by the supergravity corrections.

The relevant branch in the field space is represented by the  $SU(6)$   $D$ - and  $F$ -flat trajectory parameterized by the invariant  $\text{Tr}\Sigma^2$ . This corresponds to an arbitrary expectation value along the component

$$\Sigma = \text{diag}(1, 1, 1, -1, -1, -1)\frac{S}{\sqrt{6}}. \quad (6)$$

The key point here is that above component has no self-interaction (*i.e.*  $\text{Tr}\Sigma^3 = 0$ ) and appears in the superpotential linearly. At the generic point of this moduli space the gauge  $SU(6)$  symmetry is broken to  $SU(3) \otimes SU(3) \otimes U(1)$ . All gauge-non singlet Higgs fields are getting masses  $\mathcal{O}(S)$  and therefore, for large values of  $S$ ,  $S \gg \sqrt{\xi}$ , they decouple. Part of them gets eaten up by the massive gauge superfields. These are the components of  $\Sigma$  transforming as  $(3, \bar{3})$  and  $(\bar{3}, 3)$  under the unbroken subgroup. All other Higgs fields get large masses from the superpotential. The massless degrees of freedom along the branch are therefore : two singlets  $S$  and  $X$ , the massless  $SU(3) \otimes SU(3) \otimes U(1)$  super- Yang-Mills multiplet and the massless matter superfields.

By integrating out the heavy superfields, we can write down an effective low energy superpotential by simply using holomorphy and symmetry arguments [20]. This superpotential, as well as all gauge  $SU(6)$   $D$ -terms, is vanishing. Were not for the  $U(1)_\xi$ -gauge

symmetry, the branch parameterized by  $S$ , would simply correspond to a SUSY preserving flat vacuum direction remaining flat to all orders in perturbation theory. The  $D$ -term, however, lifts this flat direction, taking an asymptotically constant value for arbitrarily large  $S$  at the tree-level. This is because all Higgs fields with charges opposite to  $\xi$  gain large masses and decouple, and  $\xi$  can not be compensated any more (notice that heavy fields decouple in pairs with opposite charges and therefore  $\text{Tr}\mathbf{Q}$  over the remaining low energy fields is not changed). As a result, the branch of interest is represented by two massless degrees of freedom  $X$  and  $S$  whose VEVs set the mass scale for the heavy particles, and a constant tree level vacuum energy density

$$V_{\text{tree}} = \frac{g^2}{2}\langle D^2 \rangle = \frac{g^2}{2}\xi^2. \quad (7)$$

This term is responsible for inflation.

The above result, which was based on holomorphy and symmetry arguments, can be easily rederived by explicit solution of the equations of motion along the inflationary branch. For doing this, we can explicitly minimize all  $D$ - and  $F$ - terms subject to large values of  $S$  and  $X$ . The relevant part of the potential is

$$V = |F_{H'_A}|^2 + |F_{\bar{H}'_A}|^2 + \frac{g^2}{2}D^2, \quad (8)$$

since the remaining  $F$ - and  $D$ - terms are automatically vanishing as long as all other gauge-non singlet Higgses are zero. We would need to include them only if the minima of the potential (8) (subject to  $S, X \gg \xi$ ) were incompatible with such an assumption. However for the branch of our interest this turns out to be not the case.

It is easy now to check that for

$$\text{Min} \left( \left| M + aX \pm \alpha \frac{S}{\sqrt{6}} \right|, \left| M' + a'X \pm \alpha' \frac{S}{\sqrt{6}} \right| \right) > g\sqrt{\xi} \quad (9)$$

all other VEVs vanish and, therefore, a nonzero contribution to the potential comes purely from the constant  $U(1)_\xi$   $D$ -term. This is when inflation takes place: starting from some chaotic initial values of  $S$  and  $X$  for which the condition (9) is far from being satisfying, the system will slowly evolve and inflate. In each case the inflaton field is represented by the appropriate combination of  $S$  and  $X$  fields.

Whenever the condition (9) is violated, some of the  $H, \bar{H}$  components become tachionic and compensate the  $D$ -term. The system very rapidly relaxes to the supersymmetric vacuum (4) and oscillates about it. Inflation is therefore terminated by this rapid water-fall [16] and the universe undergoes a short period of reheating after which it is filled up by particles in thermal equilibrium.

As we have seen, the tree-level potential along the inflationary branch is exactly flat. Radiative corrections [4,9], however, create a logarithmic slope that drives inflaton toward the minimum (4). The origin of this correction can be understood in the following way. As we have shown, the  $S$  and  $X$  VEVs set the mass scale for the heavy particles along the inflationary branch. Thus, we can think of the low energy theories at the different points of this branch as of the same theory at the different energy scales. The gauge coupling in (7) should be understood as the running gauge coupling. This is simply due to the gauge field wave function renormalization through the loops with  $U(1)_\xi$  - charged particles  $H, \bar{H}, H', \bar{H}'$ . Since their mass is set by  $S$  and  $X$  VEVs, the nontrivial dependence on these VEVs arises, providing effective one-loop potential for the inflaton field. For large field strengths or, in other words, masses of the particles in the loop much larger than  $\sqrt{\xi}$ , this potential assumes the following form [4], [9]

$$V_{\text{inf}} = \frac{g^2}{2} \xi^2 \left( 1 + \frac{3g^2}{\pi^2} \ln \left( \left| \pm \alpha S / \sqrt{6} + aX + M \right| \left| \pm \alpha S / \sqrt{6} + a'X + M' \right| \right) \right) \quad (10)$$

This is simply the asymptotic form for  $S, X \gg \sqrt{\xi}$  of the one-loop corrected effective potential

$$V_{\text{one-loop}} = \text{Tr} (-1)^F \mathcal{M}^4 \ln \mathcal{M}^2. \quad (11)$$

The contribution to (11) comes purely from the  $H, \bar{H}', H', \bar{H}$  superfields. These are the fragments  $(1, 3), (1, \bar{3})$  and  $(3, 1), (\bar{3}, 1)$  of the  $H, \bar{H}'$  with supersymmetric masses

$$\pm \alpha S / \sqrt{6} + aX + M, \quad (12)$$

and the analogous fragments of the  $H', \bar{H}$  with supersymmetric masses

$$\pm \alpha' S / \sqrt{6} + a'X + M', \quad (13)$$

respectively. All these superfields suffer from the tree level non-supersymmetric contribution to the scalar masses from the  $U(1)_\xi$   $D$ -term equal to

$$\pm g^2 \xi, \quad (14)$$

where the sign corresponds to the  $U(1)_\xi$ -charge. All other states either have no mass-splitting due to a vanishing charge (these are  $X, \Sigma$  and the gauge fields) or have no inflaton dependent mass (these are matter fields).

As we have seen, the tree level inflationary branch is a two dimensional complex plane subject to the constraint  $S, X \gg \sqrt{\xi}$ . Classically, any path parameterized by an arbitrary combination of  $S$  and  $X$  on this manifold is exactly flat and can lead to inflation with a nearly equal chance. So classically inflation can end only when condition (9) breaks down, signaling that some of the fields become tachionic and system relaxes to the global minimum. However, as we have argued, the quantum corrections provide a slope for the inflaton field and inflation in reality may end much before the instability occurs, simply because of the breakdown of the slow roll conditions.

Let us denote the direction along which inflation is taking place by  $\phi$  and write symbolically the potential (10) as  $V_{\text{inf}} \simeq V_0(1 + c g^2 \log \phi)$ , where  $c = \frac{6}{\pi^2}$ . During the slow-roll phase, when the inflaton is rolling down from large values, the cosmic scale factor may grow by  $N$  e-foldings:

$$N \simeq \frac{8\pi}{M_{Pl}^2} \int_{\phi_N}^{\phi_e} \frac{V_0}{V'} = \frac{4\pi}{M_{Pl}^2} \frac{\phi_N^2}{g^2 c}, \quad (15)$$

where  $\phi_e$  denotes the value of the field when inflation ends. Successful inflation requires  $N \simeq 60$ .

Fluctuations arise due to quantum fluctuations in the inflaton field. We may then compute the power spectrum of quantum fluctuations, which is the Fourier transform of the two-point density autocorrelation function. It has the primordial form  $P(k) \propto k^n$ , where  $k$  is the amplitude of the Fourier wavevector and  $n$

denotes the spectral index. The measurement of the quadrupole anisotropy in the cosmic microwave background radiation detected by COBE [21] allows us to fix the parameters of the model:

$$\begin{aligned}
\left(\frac{\Delta T}{T}\right) &= \sqrt{\frac{32\pi}{45}} \frac{V_0^{3/2}}{V'(\phi_N) M_{Pl}^3} \\
&\simeq 0.3 \sqrt{\frac{N}{c}} \left(\frac{\xi}{M_{Pl}^2}\right).
\end{aligned}
\tag{16}$$

Imposing  $\left(\frac{\Delta T}{T}\right) \simeq 6 \times 10^{-6}$ , for  $c \sim \frac{6}{\pi^2}$  we get  $\sqrt{\xi} \sim 10^{16}$  GeV, which is close to the GUT scale. The spectral index is practically indistinguishable from unity,  $n-1 \simeq 1 - \frac{1}{N} \simeq 0.98$ . This recovers prediction of the scenario [4], the difference is that, since our inflation is  $D$ -dominated, we do not need any assumption about the non-minimal (quartic) terms in the Kähler potential  $(\phi^* \phi)^2$ . They do not contribute in the curvature, since  $F_\phi$  is vanishing during inflation. On the contrary, in the  $F$ -dominated scenario [4] the predictions are sensitive to the precise structure of this term [22].

One may ask whether the value of  $\sqrt{\xi}$  required by density perturbations can be motivated by realistic string theory. At this point uncertainties come from the fact that in our approximation we were treating  $\xi$  as constant (up to a course-graining scale dependence through the gauge-coupling). This is certainly justified in the effective field theory approach in which  $\xi$  is treated as an input parameter. In string theories the gauge and gravitational coupling constants are set through the expectation value of the dilaton field  $s$  and the Fayet-Iliopoulos  $D$ -term actually is a function of  $(s + \bar{s})$ . Since the dilaton potential most likely is strongly influenced by the inflationary dynamics, the actual value of  $\xi$  at the moment when observationally interesting scales crossed the horizon during inflation might be quite different from the one "observed" today. It seems that entire question is related to the problem of the dilaton stabilization and it is hard to make any definite statement without knowing the details of the dilaton dynamics during inflation. All our estimates made above are valid within an effective field theory description, in which the gauge and gravitational constants can be treated as parameters whose inflaton-dependence arises from the course-graining scale-dependence.

#### IV. THE MONOPOLE PROBLEM

In the usual hybrid inflationary scenarios [16] inflation is terminated by the rolling down of a Higgs field coupled to the inflaton and consequent phase transition with symme-

try breaking. Whenever the vacuum manifold has a non-trivial homotopy, the topological defects will form much in the same way as in the conventional thermal phase transition. Thus, the straightforward generalization of the hybrid scenario in the GUT context would result in the post-inflationary formation of the unwanted magnetic monopoles. In our scenario this disaster never happens, since the inflaton field is the GUT Higgs itself. The GUT symmetry is broken both during and after inflation and the monopoles (even if present at the early stages) get inevitably inflated away. The unbroken symmetry group along the inflationary branch is  $G_{\text{inf}} = SU(3) \otimes SU(3) \otimes U(1) \otimes SU(2) \otimes U(1)_\xi$ <sup>6</sup> which gets broken to  $G_{\text{postinf}} = SU(3) \otimes SU(2) \otimes U(1) \otimes U(1)$  modulo the electroweak phase transition (extra  $U(1)$  -factor is global). Since  $\pi_2(G_{\text{inf}}/G_{\text{postinf}}) = 0$  no monopoles are formed.

In conclusion, we have shown that a successful model of inflation may naturally arise from concrete supersymmetric Grand Unified Theories where the doublet-triplet splitting problem is solved. To achieve that, no price of enlarging the scalar sector is to be paid. The Fayet-Iliopoulos  $D$ -term of a gauge  $U(1)_\xi$  symmetry plays a crucial role both in the solution of the doublet-triplet splitting problem and in providing a suitable slope for the inflaton potential which is protected from dangerous supergravity corrections. Since the inflaton is a GUT adjoint Higgs field, the Grand Unified symmetry is broken both during and after inflation. As a result, the universe popping out after inflation is safe from monopoles.

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<sup>6</sup>If the gauge  $U(1)_\xi$  is a stringy anomalous  $U(1)$ , it will be broken by the dilaton even if all other charged fields vanish. In this case the unbroken symmetry has to be understood as a global one.

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