A Predictive Inflationary Scenario Without
The Gauge Singlet

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Abstract

We propose a new realization of the chaotic inflationary scenario
in which the scalar field responsible for inflation also spontaneously
breaks the underlying gauge symmetry at a superheavy scale \( \sim 10^{15} - 10^{17} \) GeV. A possible framework is provided by the superstring in-
spired gauge models, in which case several predictions are essentially
model independent. The spectral index for the scalar perturbations
\( n \approx 0.92 - 0.88 \), while the ratio of the tensor to the scalar quadrupole
anisotropy is \( (\Delta T/T)^2/(\Delta T/T)^{2}_S \approx 0.4 - 0.7 \). On smaller angular
scales, \( \delta T/T(1^\circ) \approx (9 - 16) \times 10^{-6} \) and \( \delta T/T(2.1^\circ) \approx (6 - 10) \times 10^{-6} \).
Implications for magnetic monopoles and cosmic strings as well as the
gauge hierarchy problem are pointed out.

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The simplest realizations of the new or the chaotic inflationary scenario\(^{(1)}\) invoke a weakly coupled scalar field which typically is a singlet under the full gauge symmetry of the model.\(^{(2,3)}\) Gauge singlets arise naturally within the framework of higher dimensional cosmologies and inflationary scenarios based on these Kaluza-Klein type ideas have been discussed quite some time ago.\(^{(4)}\) For almost a decade, however, the most promising approach for unifying grand unification with gravity has been based on superstring theories. A variety of related ideas have been explored in the literature. One of the most elegant is also the earliest,\(^{(5)}\) based on the compactification of six of the ten dimensions of the heterotic \(E_8 \times E_8\) superstring theory\(^{(6)}\) on a Calabi-Yau (C-Y) manifold. Four dimensional gauge models obtained from the C-Y approach are in many ways more satisfying than the standard supersymmetric grand unified theories (SUSY GUTS). For instance, the number of chiral families is related to the Euler character of the underlying C-Y manifold.\(^{(5)}\) Discrete symmetries such as matter parity arise automatically, and are needed to adequately suppress proton decay, ensure the existence of a pair of light higgs doublets, etc.

The main purpose of this paper is to point out an intriguing possibility, that the (chaotic) inflationary scenario perhaps could be implemented in superstring models without invoking a gauge singlet field!\(^{(7)}\) The idea is that inflation could be driven by the very same field which also is responsible for breaking the underlying gauge symmetry at some superheavy scale \(\sim 10^{15} - 10^{17} \text{ GeV}\). The presence of supersymmetry, broken at a scale \(\sim 10^{3\pm1} \text{ GeV}\), is essential to ensuring the appearance of a ‘sufficiently flat’ potential, needed both for phenomenology as well as for inflation. Moreover, the dimensionless coupling which enters in the determination of \(\delta \rho/\rho\) is related to non-perturbative (instanton) effects of the underlying string theory, thereby providing (at least) another rationale as to why it happens to be much smaller than unity. The form of the inflationary potential in this class of models is narrowly constrained, thereby allowing for several essentially model independent predictions.

In the C-Y approach to the heterotic \(E_8 \times E_8\) superstring compactification,\(^{(5)}\) the ten-dimensional spacetime splits into \(M_4 \times K\), where \(M_4\) denotes the Minkowski spacetime and \(K\) is a six-dimensional C-Y manifold. By embedding the spin connection of \(K\) in the first \(E_8\), one obtains a four-dimensional gauge model which possesses an \(E_6\) gauge symmetry and \(N = 1\) supersymmetry. The C-Y space \(K\) is usually constructed as \(K = K_0/G\), where \(G\) is
a freely acting discrete group on the simply connected C-Y space $K_0$. The number of chiral fermion families of $E_6$ is given by $|\chi(K_0)| / 2n(G)$, where $\chi(K_0)$ is the Euler character of $K_0$ and $n(G)$ is the order of $G$. Non-trivial Wilson loops on $K$ will break $E_6$ to some subgroup $H$ of rank five or six. Note that in the compactification schemes under consideration the gauge coupling typically gets related to the vacuum expectation value of some dilaton field. Precisely how this occurs, especially in the context of the early universe is an issue beyond the scope of this paper.

Consider, for definiteness, the case where $H$ is the maximal subgroup $SU(3)_c \times SU(3)_L \times SU(3)_R$ of $E_6$. The left-handed lepton, quark and anti-quark fields from the $27$ of $E_6$ transform under $H$ as

\[
\begin{align*}
\ell &= (1, 3, 3) = \begin{pmatrix} H^a & H^d & L \\ E^c & \nu^c & N \end{pmatrix} \\
q &= (3, 3, 1) = \begin{pmatrix} u \\ d \\ g \end{pmatrix} \\
Q &= (\bar{3}, 1, 3) = \begin{pmatrix} u^c \\ d^c \\ g^c \end{pmatrix}
\end{align*}
\]

After the flux breaking, in addition to the three (chiral) $27$’s, there should survive at least two massless $\ell$, $\bar{\ell}$ pairs, to provide the necessary Higgs fields $N, \bar{N}$ and $\nu^c, \bar{\nu}^c$ for the symmetry breaking of $(SU(3))^3$ to the minimal supersymmetric standard model.

A variety of observational constraints such as the suppression of proton decay, $\sin^2 \theta_W \simeq 0.23$, etc. require that the breaking of $SU(3)^3$ takes place at some superheavy scale. In order to generate this scale, the potential for the fields that acquire vevs of this order must possess $D$ and $F$ flat directions.\(^{8,9,10}\) It turns out that the $\ell^3$ and $\bar{\ell}^3$ terms (from $(27)^3$ and $(\bar{27})^3$) in the superpotential are automatically $F$-flat in the $N, \bar{N}, \nu^c, \bar{\nu}^c$ directions. To ensure the vanishing of the $D$-terms, pairs of $\ell, \bar{\ell}$ must acquire vevs along the $N, \bar{N}$ and $\nu^c, \bar{\nu}^c$ directions such that

\[
< \sum_i 27_i^a T^a 27_i > = < \sum_i \bar{27}_i^a T^a \bar{27}_i >
\]

This ensures the cancellation of quartic contributions to the potential from
the $D$-terms.

The quartic (leading non-renormalizable) contribution to the superpotential $W$ takes the generic form

$$\frac{\lambda}{M_P}(27\overline{27})^2$$

(3)

where $M_P \approx 1.2 \times 10^{19} GeV$ is the Planck scale, and $\lambda$ is a dimensionless parameter. It has been pointed out$^{(11)}$ that there should be $F$-flat directions along which the non-renormalizable contribution in (3) is generated only through the non-perturbative world-sheet instanton effects. The coefficient $\lambda$ is then proportional to $\exp(-c/g^2)(c > 0$ and $g$ denotes the world sheet coupling) and could reasonably be expected to be much smaller than unity. We will be more precise about the value of $\lambda$ when we discuss the inflationary aspects of the model.

Next we make the standard assumption that the symmetry breaking of $SU(3)^3$ to the MSSM has a radiative origin. This requires that the superpotential contains cubic couplings of $N, \bar{N}(\nu^c, \bar{\nu}^c)$ that are sufficiently strong. The coupling $gg^cN(\bar{g}\bar{g}^c\bar{N})$ is one such example. The presence of these couplings will ensure that the loop corrections will drive the mass squared term for the $N, \bar{N}(\nu^c, \bar{\nu}^c)$ pair, arising from supersymmetry breaking, to the negative values needed for the spontaneous symmetry breaking. To simplify, we henceforth base our discussion on a pair of scalar fields. [For instance, $\phi, \bar{\phi}$ could be the pair $N, \bar{N}$ which breaks $SU(3)^3$ to $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$.] In the D-flat direction the scalar potential $V(\phi)$ has the form

$$V(\phi) \approx -M_S^2 |\phi|^2 + \frac{\lambda^2}{3} \frac{|\phi|^6}{M_P^2}$$

(4)

where $M_S (\sim 10^{3\pm1} GeV)$ denotes the supersymmetry breaking scale and we assume that the coefficients of the higher order terms are small enough to make them negligible during the relevant last stages of the inflationary phase. [It remains to be seen whether this important assumption can be realized in realistic ‘string derived’ models.] Minimization of $V(\phi)$ gives

$$|<\phi>| = |<\bar{\phi}>| \equiv M \simeq \lambda^{-\frac{1}{2}}(M_PM_S)^{\frac{1}{2}} GeV$$

(5)
For values of $\phi$ larger than $M$, the $\phi^6$ term in (4) dominates. Provided that $\lambda$ is sufficiently small, this kind of potential will yield the chaotic inflationary scenario.\(^{(1)}\) The inflationary phase takes place for $\phi \gg M_P$ (with the constraint that $\lambda^2 |\phi|^6 / 3M_P^2 \lesssim M_P^2$), and ends when $\phi$ becomes of order $0.5M_P$ (for the $\phi^6$ potential). The field $\phi$ then rolls down to the minimum at $\phi = M$ and performs damped oscillations of frequency $\sim M_S$.

An estimate of the order of magnitude of $\lambda$ is obtained by considering the contribution of the scalar metric perturbation to the microwave background quadrupole anisotropy (scalar Sachs-Wolfe effect) and comparing it with the recent COBE measurement.\(^{(12)}\) One has\(^{(1)}\) (the subscript $S$ denotes the scalar contribution):

$$\left(\frac{\Delta T}{T}\right)^2_S \simeq \frac{32\pi}{45} \frac{V^3}{V''M_P^6}|_{k \sim H}$$

where the right hand side is to be evaluated when the scale $k^{-1}$, corresponding to the present horizon size, crossed inside the horizon during inflation. Equation (6) can be re-written as

$$\left| \left(\frac{\Delta T}{T}\right)_S \right| \simeq 0.023\lambda N_H^2$$

where $N_H = \frac{2\pi}{3}(\phi/M_P)^2|_{k \sim H}$ denotes the corresponding number of e-foldings. Taking $N_H$ on the order of 50, [this is somewhat smaller than the usually quoted value of 60 due to the lower damping rate of the oscillating $\phi$ field], and $\Delta T/T \approx 6 \times 10^{-6}$ from COBE, we estimate the fundamental quantity $\lambda$ to be of order $10^{-7}$.

Inserting this value of $\lambda$ in (5), we see that the vev $M \simeq 10^{14.5}$ $(M_S/10^3 \ GeV)^{\frac{1}{2}} \ GeV$. In order to estimate the decay width of $\phi$ one needs to know the relevant couplings. Clearly, since the decay products have masses $\lesssim M_S$ ($\approx M_0$), these couplings all arise from the non-renormalizable terms (with suppressed couplings) in the superpotential. Some typical ones are $HH\phi^2/M_P^4$, $\nu^c\nu^c\phi^2/M_P^4$, etc. The first one produces higgsinos as decay products, while the second coupling gives rise to ‘right handed’ neutrinos. The decay width $\Gamma$ of $\phi$ is (roughly) estimated to be

$$\Gamma \sim \mathcal{O}(10^{-1})(M_S^3/M^2)$$

$$\sim \mathcal{O}(10^{-21})(M_S/10^3 \ GeV)^2 GeV \quad (8)$$
The oscillations of $\phi$ are damped out when the Hubble time $t$ becomes $\sim \Gamma^{-1}$, and the universe ‘reheats’ to a temperature

$$T_r \sim (\Gamma M_P)^{1/2} \sim 10^{-1}(M_S/10^3 \text{ GeV})\text{GeV} \quad (9)$$

An inflationary scenario is certainly incomplete without an explanation of the origin of the observed baryon asymmetry in the universe. This is a particularly pressing issue in the present case. The reheat temperature is quite low ($\lesssim \text{few GeV}$), so that some of the more interesting (from the inflationary viewpoint) scenarios, such as baryons from leptons\textsuperscript{(13)} or electroweak baryogenesis,\textsuperscript{(14)} are not applicable. Actually, the problem has been discussed in some detail in an earlier work\textsuperscript{(15)}. Here, for completeness, we present only the essential idea, keeping details to a minimum. The baryon asymmetry is given by the formula

$$n_b/s \sim \frac{T_r \Gamma_{\Delta B \neq 0}}{M_\phi} \quad (10)$$

where $\Gamma(\Gamma_{\Delta B \neq 0})$ denotes the total (baryon number violating) decay width of $\phi$. Consider the superpotential couplings $gg^c\phi$, $g^uw^cd^c$ and $gd^c\nu^c$. The coefficient in front of the first coupling is assumed to be of order $M_g/\langle \phi \rangle$, while the remaining two couplings carry coefficients of order unity. The decay width for the baryon number violating process $\phi \to u^c d^c d^c \nu^c$ is then given by

$$\Gamma_{\Delta B \neq 0} \sim \frac{1}{16\pi} \left( \frac{1}{8\pi^2} \right)^3 \left( \frac{M_g}{\langle \phi \rangle} \right)^2 \frac{M_\phi^5}{M_g^4} \cdot \text{(no. of channels)} \quad (11)$$

Here $M_g$ denotes the mass of the $g$ boson and in estimating the number of channels we include the sum over color and flavors. The baryon asymmetry is estimated to be

$$n_b/s \sim \mathcal{O}(10^{-1}) \left( \frac{1}{8\pi^2} \right)^3 \frac{T_r M_\phi}{M_g^2} \quad (12)$$

A number of comments are in order:
1. The $g$ boson mass should be $\sim 10^5 - 10^6 GeV$ in order to generate $n_b/s \sim 10^{-10} - 10^{-11}$. The scenario actually requires (a minimum of) two species of $g$’s. Precise details would be model dependent.

2. The presence of such relatively ‘light’ $g$’s would impose constraints on the model arising from proton decay, $n-\bar{n}$ oscillations, etc.

3. Since the reheat is so low, 2-2 scatterings cannot wipe out the asymmetry generated above. This is certainly a plus for the model.

A second scenario for implementing baryogenesis at low ($\sim GeV$ - few $MeV$) temperature has previously been discussed in ref (16). It needs the presence in the superpotential of the baryon number violating operator $u^c d^c d^c$ which presumably is a mild requirement. This scenario also seems to fit well with the present inflationary framework.

We now turn to the important issue of topological defects. Depending on the model, magnetic monopoles and/or cosmic strings will arise through the Kibble mechanism(17) at the end of inflation. For instance, the breaking of $SU(3)^3$ produces magnetic monopoles. Remarkably, however, there is no monopole problem.(18) Two crucial differences from ordinary GUTS are i) the higgs correlation length $\xi$ is of order $M^{-1}$ and not $M^{-1}$, and ii) $\phi$ dominates for quite some time after the production of the topological defects.

Monopoles are produced via the Kibble mechanism when the $\phi$-field oscillations over the barrier at $\phi = 0$ with height $M_S^2 M^2$ come to a halt. Their initial number density is

$$n_M \sim \frac{p}{\frac{4}{3} \xi M} \sim 10^{-2} M_S^3$$

where $p \sim 10^{-1}$ is a geometric factor. Consequently, the initial monopole energy density is given by

$$\frac{\rho_M}{\rho_\phi} \sim 10^{-2} \frac{M_S^3 m_M}{M_S^4 \langle \phi \rangle^2} \sim 10^{-2} \frac{M_S m_M}{\langle \phi \rangle^2}$$

where $m_M$ denotes the monopole mass. The ratio in (14) remains constant until radiation takes over at $T_r$. Assuming that no further entropy is generated, one finds that $r \equiv n_M/s \sim 10^{-2} \frac{M_S T_r}{\langle \phi \rangle^2} \sim 10^{-20}$ for $M_S \sim TeV$, $\langle \phi \rangle \sim 10^{15} GeV$. 


We therefore conclude that if the inflationary scenario is implemented within an \((SU(3))^3\) model, one expects to see magnetic monopoles (carrying three quanta of Dirac magnetic charge) at or close to the Parker bound!

Depending on the model, cosmic strings can be produced (analogous to the magnetic monopoles) at the end of the inflationary epoch. Their thickness is of order \(M_S^{-1}\), while their mass per unit length is of order \(M^2\). One needs \(M \sim 10^{16} \text{ GeV}\) for strings to play a significant role in large scale structure formation. This can be achieved by choosing \(\lambda\) to be somewhat smaller than \(10^{-7}\), in which case the main source of primordial density fluctuations would be due to cosmic strings.

Before proceeding to a discussion of some model independent predictions of this inflationary scenario, we wish to go back to the superpotential in this class of models. With some clever symmetries it is possible, in principle, to eliminate the lowest order non-renormalizable term in (3). In this case the leading non-renormalizable term in the superpotential is \(\mathcal{O}(\frac{1}{M_P^4})\lambda'(27 \bar{27})^3\). We then expect the effective potential \(V(\phi)\) to have the form (assuming that the higher order terms can be ignored; see remarks immediately preceding eq. (5)):

\[
V(\phi) \approx -M_S^2 |\phi|^2 + \frac{\lambda'^2}{5} \frac{|\phi|^{10}}{M_P^6}
\]  

Minimization then yields the vev to be

\[
|<\phi>| \equiv |<\bar{\phi}>| \equiv M' \simeq \lambda'^{-\frac{1}{2}} (M_S^2 M_P^6)^{\frac{1}{2}} \text{GeV}
\]  

The quantity \((\frac{\Delta T}{T})_S\) in this case is proportional to \(N_k^3\). Proceeding as in the previous case, one finds that the dimensionless parameter \(\lambda' \approx 0.2 \times 10^{-8}\). Substitution in (16) yields

\[
|<\phi>| \equiv |<\bar{\phi}>| = M' \simeq 10^{17} \left(\frac{M_S}{10^3 \text{ GeV}}\right)^{\frac{1}{4}} \text{ GeV}
\]  

The scale in (17) is somewhat larger than the typical SUSY GUT scale of \(10^{16} \text{ GeV}\), although this need not be an issue. However, the ‘reheat’ temperature is in the MeV range at best, and so it should be clear that in order to have the standard nucleosynthesis scenario the leading non-renormalizable terms in the superpotential should not be of dimension higher than seven.
The effective potential during the inflationary phase therefore will be assumed to be proportional either to $\phi^6$ or $\phi^{10}$. Note that in the latter case the topological defects become less interesting. The monopole number density will be extremely small (see eq. (14)), while the cosmic strings are excessively massive.

In chaotic inflation with a $\phi^7$ scalar potential, the amplitude of the density perturbation on a given scale $k^{-1}$, as it crosses inside the horizon is proportional to $N_k^{2+2} \approx N_H^{2+2} (k^{-1}(\text{Mpc})/10^4)^{2+2}$. Taking $N_H \approx 50$, this implies that the spectral index $n \approx 0.92(0.88)$ for $\gamma = 6(10)$. Recall that $n = 1$ corresponds to the Harrison-Zeldovich case.

Employing the well known relation for the gravitational wave contribution to the quadrupole anisotropy\(^{(19)}\)

$$\left(\frac{\Delta T}{T}\right)_T^2 \approx 0.6 \frac{V}{M_P}$$

we find

$$r \equiv \frac{(\Delta T/T)_P^2}{(\Delta T/T)_S^2} \approx \frac{3.4n}{N_H}$$

$$\approx 0.4(0.7) \text{ for } \gamma = 6(10)$$

Knowing $n$ and $r$ we can estimate the bias factor $b(\equiv \sigma_8^{-1}$, where $\sigma_8$ is the rms mass fluctuation on the scale $8h^{-1}\text{ Mpc}$) for a cold dark matter scenario using the approximate relation\(^{(20)}\)

$$b_{CDM} \approx 100^{(1-n)/2} \sqrt{1 + r} \approx 1.4 - 1.7$$

For a mixed (cold +20% hot) dark matter scenario,\(^{(21,22)}\) the bias factor turns out to be $b_{MDM} \approx 1.5b_{CDM}$.

Our final topic concerns the anisotropy predictions on angular scales of $1^\circ$ and $2.1^\circ$. We will follow ref. (23), taking into account the following. Firstly, it has been noted\(^{(24)}\) that the COBE DMR gives $Q_{rms-PS} \approx 14\mu K \pm 27\%$, which corresponds to a reduction of the published COBE numbers by $\approx 15\%$. Secondly, the power spectrum here has less power on smaller scales and in addition, the tensor contribution to the quadrupole anisotropy is not negligible. Taking all this into account we find that, unless reionization was important, $\delta T/T(1^\circ) \approx (9 - 16) \times 10^{-6}$ and $\delta T/T(2.1^\circ) \approx (6 - 10) \times 10^{-6}$. 

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To conclude, the proposal outlined above for implementing inflation could, also, in principle, be considered within the framework of ordinary supersymmetric GUTS. One would have to ensure, through suitable symmetries, that the inflationary potential is consistent with all of the phenomenological constraints. The superstring framework (Calabi-Yau, orbifolds, four dimensional constructions,...) appears, however, to provide a more natural framework. The value of the dimensionless coupling $\lambda$ (or $\lambda'$) [and also presumably of other couplings associated with the leading non-renormalizable terms in the superpotential], is determined to be $\sim 10^{-7} - 10^{-8}$, which is precisely what one needs to ensure the existence of a pair of ‘light’ ($\sim M_S$) higgs doublets (assuming, of course, that the doublets acquire their mass only through the quartic non-renormalizable couplings). Consequently, the doublets should be protected from acquiring large masses through cubic couplings in order to resolve the gauge hierarchy problem. Finally, we have concentrated in this work on outlining the scenario and describing some model independent predictions. It would be extremely interesting to find realistic examples of models in which the coupling $\lambda$ turns out to be of the right order of magnitude.

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