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No scale SUGRA SO(10) derived Starobinsky model of inflation

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ABSTRACT

We show that a supersymmetric renormalizable theory based on gauge group SO(10) and Higgs system $\mathbf{10} \oplus \mathbf{210} \oplus \mathbf{126} \oplus \overline{\mathbf{126}}$ with no scale supergravity can lead to a Starobinsky kind of potential for inflation. Successful inflation is possible in the cases where the potential during inflation corresponds to $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, $SU(5) \times U(1)$ and flipped $SU(5) \times U(1)$ intermediate symmetry with a suitable choice of superpotential parameters. The reheating in such a scenario can occur via non-perturbative decay of inflaton i.e. through "preheating". After the end of reheating, when universe cools down, the finite temperature potential can have a minimum which corresponds to MSSM.

1. Introduction

The theory of cosmological inflation [1–3] not only solves the problems (flatness, horizon, etc.) of standard big bang theory, but also explains the seed fluctuations which can grow via gravitational instability to form the large scale structure of the universe [4]. There are stringent constraints on inflationary theories from CMB observations [5–8] and many of the generic models like the quartic potential and quadratic potential are either ruled out or disfavored by the bound on the tensor to scalar ratio which is $r_{0.05} < 0.12$ at 95% CL from joint analysis of BICEP2/Keck array and Planck data [9]. Among the generic inflation models which survive the stringent constraint on r is the R^2 inflation model of Starobinsky [1] which predicts $n_s - 1 = -2/N$ and $r = 12/N^2 \sim 0.002$–0.004. The theoretical motivation for the Starobinsky model is provided in [10] where it has been shown that the Starobinsky potential for inflation can be derived from supergravity (SUGRA) with a no-scale [11–13] Kähler potential and a Wess Zumino superpotential with specific couplings. Supergravity models of inflation based on the Jordan frame supergravity [14–16] and D-term superpotential [17] also give inflationary potential which is identical to the Starobinsky potential at large field values. The natural choice for the inflaton in supergravity models is the Higgs fields of the grand unified theories. A no-scale SUGRA model of inflation based on the SU(5) GUT using the **24**, **5** and $\bar{\mathbf{5}}$ Higgs in the superpotential has been constructed [18]. The SU(5) symmetry breaks to MSSM with the appropriate choice of *vev* for the **24** and a D-flat linear combination of H_u and H_d of MSSM acts as the inflaton [18].

In the present work we study inflation in a renormalizable grand unified theory based on the SO(10) gauge group with no scale SUGRA. Inflation in the context of SUSY SO(10) has been studied earlier in [19–23] with the SO(10) invariant superpotential with the minimal Kähler potential which gives polynomial potentials of inflation. In this paper we show that a renormalizable Wess–Zumino superpotential of SO(10) GUT along with no-scale Kähler potential can give us Starobinsky kind of inflationary potential with specific choice of superpotential parameters. The Higgs supermultiplets we consider are **10**, **210**, **126** ($\overline{\mathbf{126}}$). Among these, the **210** and **126** ($\overline{\mathbf{126}}$) are responsible for breaking of SO(10) symmetry down to MSSM. The **210** supermultiplet alone can give different intermediate symmetries [24] depending upon which of its MSSM singlet fields takes a *vev*. Then **126** ($\overline{\mathbf{126}}$) breaks this intermediate symmetry to MSSM. We find that successful inflationary potential can be achieved in the case of $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, $SU(5) \times U(1)$ and flipped $SU(5) \times U(1)$ symmetry. The other possible intermediate symmetries of Pati–Salam ($SU(4)_C \times SU(2)_L \times SU(2)_R$) or $SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B-L}$ gauge groups do not give phenomenologically correct inflationary potentials.

At the end of inflation, the reheating can occur via non-perturbative decay of inflaton to bosons of the intermediate scale model. After the end of reheating, when universe cools down, the

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finite temperature potential can have a minimum which corresponds to MSSM and the universe rolls down to this minimum at temperature $<< T_R$ (reheat temperature).

2. Inflation in SO(10) with no scale SUGRA

The minimal supersymmetric grand unified theory based on SO(10) gauge group [24–28] has $\mathbf{10}(H_i)$, $\mathbf{210}(\Phi_{ijkl})$ and $\mathbf{126}(\Sigma_{ijklm})$ ($\overline{\mathbf{126}}(\overline{\Sigma}_{ijklm})$) Higgs supermultiplets. The representations: H_i is 1 index real, Σ_{ijklm} is complex (5 index, totally-antisymmetric, self-dual) and Φ_{ijkl} is 4 index totally-antisymmetric tensor. Here $i, j, k, l, m = 1, 2 \ldots 10$ run over the vector representation of SO(10). The renormalizable superpotential for the above mentioned fields is given by

$$W = \frac{m_\Phi}{4!}\Phi^2 + \frac{\lambda}{4!}\Phi^3 + \frac{m_\Sigma}{5!}\Sigma\overline{\Sigma} + \frac{\eta}{4!}\Phi\Sigma\overline{\Sigma} + m_H H^2 + \frac{1}{4!}\Phi H(\gamma\Sigma + \bar{\gamma}\overline{\Sigma}). \quad (1)$$

The no-scale form of Kähler potential is taken to be

$$K = -3\ln(T + T^* - \frac{1}{3}(\frac{1}{4!}\Phi^\dagger\Phi + \frac{1}{5!}\Sigma^\dagger\Sigma + \frac{1}{5!}\overline{\Sigma}^\dagger\overline{\Sigma} + H^\dagger H)). \quad (2)$$

Here T is the single modulus field arising due to string compactification and we are taking $M_P = 1$.

The **10** and $\overline{\mathbf{126}}$ are required for Yukawa terms to give masses to the fermions while **126** ($\overline{\mathbf{126}}$) breaks the SO(10) gauge symmetry to MSSM together with **210**-plet. However to have an intermediate symmetry rather than MSSM, the **210**-plet Higgs is sufficient. It can lead to various possible intermediate symmetries depending on which components of the **210**-plet take *vevs*. The decomposition of Higgs supermultiplets required for SO(10) symmetry breaking in terms of Pati–Salam gauge group ($SU(4)_C \times SU(2)_L \times SU(2)_R$) is given by [29]

$$\begin{aligned} 210 &= (15,1,1) + (1,1,1) + (15,1,3) + (15,3,1) \\ &\quad + (6,2,2) + (10,2,2) + (\bar{10},2,2), \\ 126 &= (\bar{10},1,3) + (10,3,1) + (6,1,1) + (15,2,2), \\ \overline{126} &= (\bar{10},3,1) + (10,1,3) + (6,1,1) + (15,2,2). \end{aligned} \quad (3)$$

The field components which will not break the MSSM symmetry are allowed to take *vevs*. In this case they are [28]

$$\begin{aligned} p &= \langle\Phi(1,1,1)\rangle, \; a = \langle\Phi(15,1,1)\rangle, \\ \omega &= \langle\Phi(15,1,3)\rangle, \; \sigma = \langle\Sigma(\bar{10},3,1)\rangle, \\ \bar{\sigma} &= \langle\overline{\Sigma}(10,3,1)\rangle. \end{aligned} \quad (4)$$

The superpotential in terms of these *vevs* is

$$W = m(p^2 + 3a^2 + 6\omega^2) + 2\lambda(a^3 + 3p\omega^2 + 6a\omega^2) + m_\Sigma\sigma\bar{\sigma} + \eta\sigma\bar{\sigma}(p + 3a - 6\omega). \quad (5)$$

The vanishing of D-terms gives the condition $|\sigma| = |\bar{\sigma}|$ [28]. The symmetry breaking path of SO(10) is

$$\mathrm{SO}(10) \xrightarrow{210} \text{Intermediate symmetry} \xrightarrow{126} MSSM.$$

For the first step symmetry breaking one can set $|\sigma| = |\bar{\sigma}| = 0$. Then the possible intermediate symmetries with **210** only are [28]:

1. If $a \neq 0$ and $p = \omega = 0$, it gives $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B\text{-}L}$ symmetry.
2. If $p \neq 0$ and $a = \omega = 0$, this results in $SU(4)_C \times SU(2)_L \times SU(2)_R$ symmetry.
3. If $\omega \neq 0$ and $p = a = 0$, it gives $SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B\text{-}L}$ symmetry.
4. If $p = a = -\omega \neq 0$, this has $SU(5) \times U(1)$ symmetry.
5. If $p = a = \omega \neq 0$, $SU(5) \times U(1)$ symmetry but with flipped assignments for particles.

The superpotential in terms of *vevs* of **210** is given by

$$W = m(p^2 + 3a^2 + 6\omega^2) + 2\lambda(a^3 + 3p\omega^2 + 6a\omega^2). \quad (6)$$

Here $m = m_\Phi$. Similarly no-scale Kähler potential is

$$K = -3\ln(T + T^* - \frac{1}{3}(|p|^2 + 3|a|^2 + 6|\omega|^2)). \quad (7)$$

The F-term potential has the following form,

$$V = e^G\left[\frac{\partial G}{\partial\phi^i}K^i_{j^*}\frac{\partial G}{\partial\phi_{j^*}} - 3\right], \quad (8)$$

where

$$G = K + \ln W + \ln W^*. \quad (9)$$

The kinetic term is given as $K_i^{j^*}\partial\phi^i\partial\phi_{j^*}$. Here i runs over different fields T, p, a and ω. $K^i_{j^*}$ is the inverse of Kähler metric $K_i^{j^*}$ given by

$$K_i^{j^*} = \frac{1}{\Gamma^2}\begin{pmatrix} 3 & -p^* & -3a^* & -6\omega^* \\ -p & \Gamma + \frac{1}{3}|p|^2 & a^*p & 2\omega^*p \\ -3a & ap^* & 3\Gamma + 3|a|^2 & 6a\omega^* \\ -6\omega & 2\omega p^* & 6a^*\omega & 6\Gamma + 12|\omega|^2 \end{pmatrix}, \quad (10)$$

where $\Gamma = T + T^* - \frac{1}{3}(|p|^2 + 3|a|^2 + 6|\omega|^2)$. After simplifying, the potential given by Eq. (8) has the following form,

$$V = \frac{1}{\Gamma^2}\left|\frac{\partial W}{\partial\phi_i}\right|^2. \quad (11)$$

We assume that the non-perturbative Planck scale dynamics [18,10,30] fixes the values of $T = T^* = \frac{1}{2}$. After fixing the *vev* for T the kinetic terms of T can be neglected. We study all possible cases of intermediate symmetries mentioned earlier for inflationary conditions in SO(10) with no-scale SUGRA. For simplicity we assume our fields to be real.

Case I: $a \neq 0$ and $p = \omega = 0$, $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B\text{-}L}$ symmetry.

The kinetic and potential energy terms are given by

$$\begin{aligned} L_{K.E.} &= \frac{(1-a^2)(\partial_\mu p)^2 + 3(\partial_\mu a)^2 + 6(1-a^2)(\partial_\mu\omega)^2}{(1-a^2)^2}, \\ V &= \frac{36a^4\lambda^2 + 72a^3\lambda m + 36a^2m^2}{(1-a^2)^2}. \end{aligned} \quad (12)$$

To get the canonical K.E. terms we need to redefine our fields in terms of new fields χ_1, χ_2, χ_3,

$$a = \tanh[\frac{\chi_1}{\sqrt{3}}], \; p = \mathrm{sech}[\frac{\chi_1}{\sqrt{3}}]\chi_2, \; \omega = \frac{1}{\sqrt{6}}\mathrm{sech}[\frac{\chi_1}{\sqrt{3}}]\chi_3. \quad (13)$$

The potential $V(\chi_1, \chi_2, \chi_3)$ is flat along χ_1 direction for $\chi_2 = \chi_3 = 0$ and is confined in the orthogonal (χ_2, χ_3) directions as shown in Fig. 1.