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

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A R T I C L E   I N F O

A B S T R A C T

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.

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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

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

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

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

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

The $\mathbf{10}$ and $\overline{\mathbf{126}}$ are required for Yukawa terms to give masses to the fermions while $\mathbf{126}$ ($\overline{\mathbf{126}}$) breaks the SO(10) gauge symmetry to MSSM together with $\mathbf{210}$-plet. However to have an intermediate symmetry rather than MSSM, the $\mathbf{210}$-plet Higgs is sufficient. It can lead to various possible intermediate symmetries depending on which components of the $\mathbf{210}$-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]

$$210 = (15, 1, 1) + (1, 1, 1) + (15, 1, 3) + (15, 3, 1)$$
$$+ (6, 2, 2) + (10, 2, 2) + (\bar{10}, 2, 2),$$
$$126 = (\bar{10}, 1, 3) + (10, 3, 1) + (6, 1, 1) + (15, 2, 2),$$
$$\overline{126} = (\bar{10}, 3, 1) + (10, 1, 3) + (6, 1, 1) + (15, 2, 2). \tag{3}$$

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

$$p = \langle\Phi(1, 1, 1)\rangle, \ a = \langle\Phi(15, 1, 1)\rangle,$$
$$\omega = \langle\Phi(15, 1, 3)\rangle, \ \sigma = \langle\Sigma(\bar{10}, 3, 1)\rangle,$$
$$\bar{\sigma} = \langle\overline{\Sigma}(10, 3, 1)\rangle. \tag{4}$$

The superpotential in terms of these *vevs* is

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

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

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

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

1. If $a \neq 0$ and $p = \omega = 0$, it gives $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ 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 $SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B-L}$ 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 $\mathbf{210}$ is given by

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

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

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

The F-term potential has the following form,

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

where

$$G = K + \ln W + \ln W^*. \tag{9}$$

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

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

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,

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

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 $p = \omega = 0$, $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ symmetry.

The kinetic and potential energy terms are given by

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

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

$$a = \tanh[\frac{\chi_1}{\sqrt{3}}], \ p = \text{sech}[\frac{\chi_1}{\sqrt{3}}]\chi_2, \ \omega = \frac{1}{\sqrt{6}}\text{sech}[\frac{\chi_1}{\sqrt{3}}]\chi_3. \tag{13}$$

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.
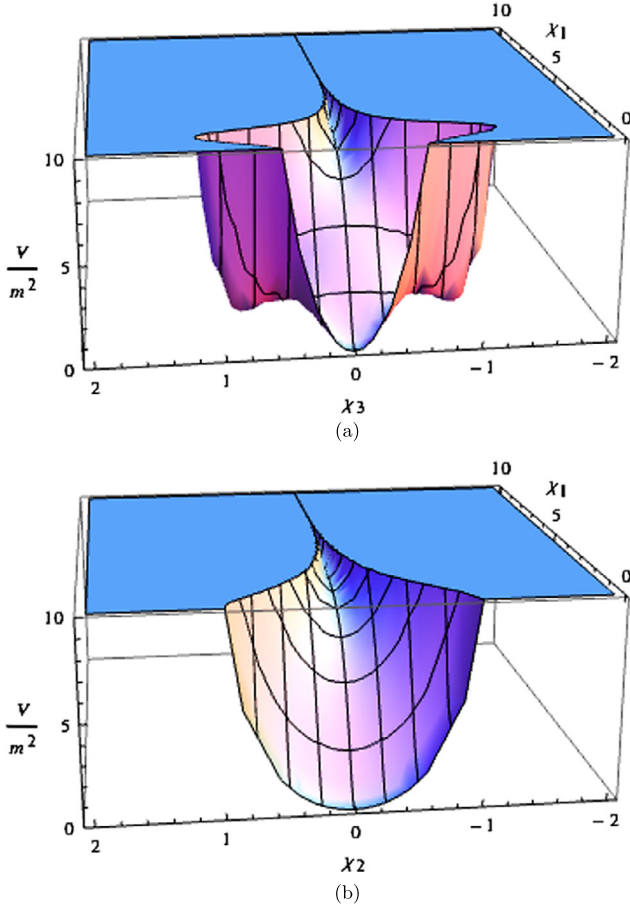
(a)

(b)

Fig. 1. The potential for the $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ intermediate symmetry is shown. The inflation potential is along χ_1 direction. In Fig. 1(a) we show $V(\chi_1, \chi_2 = 0, \chi_3)$ and in Fig. 1(b) $V(\chi_1, \chi_2, \chi_3 = 0)$. We see that potential is flat along χ_1 and confined along χ_2 and χ_3 respectively.

The potential $V(\chi_1)$ in the limit $\chi_2 = \chi_3 = 0$ is

$$V = \frac{36\lambda^2 \tanh^4\left[\frac{\chi_1}{\sqrt{3}}\right] + 72m\lambda \tanh^3\left[\frac{\chi_1}{\sqrt{3}}\right] + 36m^2 \tanh^2\left[\frac{\chi_1}{\sqrt{3}}\right]}{\left(1 - \tanh^2\left[\frac{\chi_1}{\sqrt{3}}\right]\right)^2}.$$

(14)

If we take $\lambda = -m$, this gives us the Starobinsky type of inflationary potential. The potential in this specific case is

$$V = 36m^2(1 - e^{-\frac{2\chi_1}{\sqrt{3}}})^2.$$

(15)

This potential is shown in Fig. 2 along with small deviations from the relation $\lambda = -m$. The slow roll parameters for this potential are given by

$$\eta = -\frac{8e^{\frac{-2\chi_1}{\sqrt{3}}}\left(1 - 2e^{\frac{-2\chi_1}{\sqrt{3}}}\right)}{3\left(1 - e^{-\frac{2\chi_1}{\sqrt{3}}}\right)^2}; \quad \epsilon = \frac{8e^{-\frac{4\chi_1}{\sqrt{3}}}}{3\left(1 - e^{-\frac{2\chi_1}{\sqrt{3}}}\right)^2}.$$

(16)

Inflation ends when $\eta \approx 1$, which corresponds to field value of $\chi_1^{end} \approx 0.5$. To have sufficient inflation which corresponds to $N_{e\text{-folds}} = 55$ gives the initial field value of $\chi_1 \approx 4.35$. The power spectrum for scalar perturbation P_R is

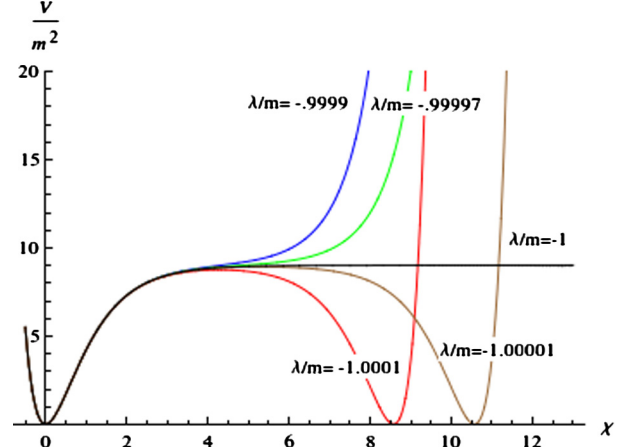$$P_R = \frac{V}{24\pi^2 \epsilon} = \frac{9m^2 \sinh^4\left(\frac{\chi_1}{\sqrt{3}}\right)}{\pi^2}.$$

(17)

Fig. 2. The potential V/m^2 for Case I for different chosen values of λ/m.

The value of $P_R = (1.610 \pm 0.01) \times 10^{-9}$ given by Planck data [7] requires value of $m = 1.311 \times 10^{-6}$ in Planck units. The spectral index $n_s = .964$ and tensor to scalar perturbation ratio $r = .002$ for $N_{e\text{-folds}} = 55$. Varying λ/m in the range (from -1.0001 to -0.9999) gives n_s in the range (0.92–1.0) and r in range (0.002–0.008).

Case II: $p \neq 0$ and $a = \omega = 0$, $SU(4)_C \times SU(2)_L \times SU(2)_R$ symmetry.

The kinetic and potential energy terms are given by

$$L_{K.E.} = \frac{(\partial_\mu p)^2 + 3(1 - \frac{p^2}{3})(\partial_\mu a)^2 + 6(1 - \frac{p^2}{3})(\partial_\mu \omega)^2}{(1 - \frac{p^2}{3})^2},$$

$$V = \frac{4m^2 p^2}{(1 - \frac{p^2}{3})^2}.$$

(18)

The fields transformations which make kinetic energy term canonical are

$$p = \sqrt{3}\tanh[\frac{\chi_1}{\sqrt{3}}], \ a = \text{sech}[\frac{\chi_1}{\sqrt{3}}]\frac{\chi_2}{\sqrt{3}}, \ \omega = \text{sech}[\frac{\chi_1}{\sqrt{3}}]\frac{\chi_3}{\sqrt{6}}. \quad (19)$$

Then the potential $V(\chi_1)$ in the limit $\chi_2 = \chi_3 = 0$ is

$$V = 3m^2 \sinh[\frac{2\chi_1}{\sqrt{3}}]^2.$$

(20)

This type of potential increases exponentially with χ_1 and is too steep to obey the slow roll conditions. The spectral index n_s has negative values over a wide range of field value and hence doesn't satisfy the inflationary constraints on scale invariance of scalar perturbations from observations.

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

The kinetic and potential energy terms are given by

$$L_{K.E.} = \frac{(1 - 2\omega^2)(\partial_\mu p)^2 + 3(1 - 2\omega^2)(\partial_\mu a)^2 + 6(\partial_\mu \omega)^2}{(1 - 2\omega^2)^2},$$

$$V = \frac{144m^2 w^2 + 180\lambda^2 w^4}{(1 - 2w^2)^2}.$$

(21)

The fields transformations which make kinetic energy term canonical are

$$\omega = \frac{1}{\sqrt{2}}\tanh[\frac{\chi_1}{\sqrt{3}}], \ p = \text{sech}[\frac{\chi_1}{\sqrt{3}}]\chi_2, \ a = \text{sech}[\frac{\chi_1}{\sqrt{3}}]\frac{\chi_3}{\sqrt{3}}. \quad (22)$$

Then the potential $V(\chi_1)$ in the limit $\chi_2 = \chi_3 = 0$ is

$$V = 72m^2 \sinh[\frac{\chi_1}{\sqrt{3}}]^2 (\cosh[\frac{\chi_1}{\sqrt{3}}]^2 + \alpha \sinh[\frac{\chi_1}{\sqrt{3}}]^2). \quad (23)$$

Here $\alpha = 5\lambda^2/8m^2$. In this case for $\alpha \geq -1$ potential increases exponentially with χ_1 and hence gives similar results as in Case II. For $\alpha < -1$ potential energy becomes negative for $\chi_1 \gtrsim 1$ and grows with large values of χ_1. Therefore this intermediate symmetry doesn't give successful inflation.

Case IV: If $p = a = \pm \omega \neq 0$, $SU(5) \times U(1)$ symmetry.

In this case we take $p = a = \pm \omega = x$, then the K.E. term and potential are given by

$$L_{K.E.} = \frac{90(\partial_\mu x)^2}{(3 - 10x^2)^2},$$

$$V = \frac{184m^2 x^2 + 1104\lambda m x^3 + 1656\lambda^2 x^4}{\left(1 - \frac{10x^2}{3}\right)^2}. \quad (24)$$

The field redefinition $x = \sqrt{\frac{3}{10}} \tanh[\frac{\chi_1}{\sqrt{3}}]$ which makes kinetic energy term canonical gives the form of potential,

$$V = 55.2m^2 (1 - e^{-\frac{2\chi_1}{\sqrt{3}}})^2, \quad (25)$$

for $\lambda = -\frac{1}{3}\sqrt{\frac{10}{3}}m$. This is a Starobinsky inflationary potential but with different relation among superpotential parameters m and λ in comparison to the Case I. In this case value of $m = 1.06 \times 10^{-6}$ is required to satisfy the constraints from CMB observations. Small variations from the relation $\lambda = -\frac{1}{3}\sqrt{\frac{10}{3}}m$ gives the same types of deviations in the Starobinsky potential as shown in Fig. 2.

At the end of inflation the inflaton χ_1 can decay to scalar bosons which have a trilinear term with Φ in superpotential e.g. $\Phi H(\gamma \Sigma + \bar{\gamma} \bar{\Sigma})$. Then the $K_\Sigma^{\Sigma^*} |W_\Sigma|^2$ and $K_{\bar{\Sigma}}^{\bar{\Sigma}^*} |W_{\bar{\Sigma}}|^2$ type of terms gives

$$V \supset ((|\gamma|^2 + |\bar{\gamma}|^2)|H|^2 + |\gamma|^2|\Sigma|^2 + |\bar{\gamma}|^2|\bar{\Sigma}|^2)|\sinh[\frac{\chi_1}{\sqrt{3}}]|^2. \quad (26)$$

Near the origin $\sinh[\frac{\chi_1}{\sqrt{3}}] \approx \frac{\chi_1}{\sqrt{3}}$, so

$$V \supset ((|\gamma|^2 + |\bar{\gamma}|^2)|H|^2 + |\gamma|^2|\Sigma|^2 + |\bar{\gamma}|^2|\bar{\Sigma}|^2)|\frac{\chi_1}{\sqrt{3}}|^2. \quad (27)$$

In our case the perturbative decay of inflaton to scalars is not efficient for typical values of $\gamma, \bar{\gamma} \sim O(.1-1.0)$ [31]. However inflaton χ_1 can decay non-perturbatively to scalar bosons leading to preheating. In [32] the mechanism of preheating in broad resonance regime has been worked out. There is another efficient way of preheating called "instant preheating" [33]. This mechanism is based upon the non-perturbative decay of inflaton to scalar bosons (in this case) when it is close to the minimum of the potential (at $\chi_1 = 0$). The particles thus produced (having mass directly proportional to the instantaneous *vev* of inflaton) decay further when inflaton rolls uphill, to the modes which are not directly coupled to inflaton. This happens because at the time of their production, their mass is zero since $\chi_1 = 0$, but as inflaton rolls back to its maximum value they become heavy so their decay width increases. In our case, every time inflaton crosses the origin it produces the H, Σ and $\bar{\Sigma}$. These decay further into the SM fermions and the right-handed neutrinos through Yukawa couplings. With this kind of chain reaction we can have an efficient way to transfer the whole energy of inflaton into relativistic particles within

few oscillations. This whole process leads to a radiation dominated universe with reheat temperature,

$$T_R \sim V_0^{1/4} \sim (m^2 \chi_1^2)^{1/4} \sim (10^{-18} M_P^4)^{1/4} \sim 10^{14} \text{ GeV}. \quad (28)$$

At the end of reheating, the universe has a finite temperature potential and after cooling from $T_R = 10^{14}$ GeV to temperature $<< T_R$, we assume that universe settles to the minimum of potential corresponding to MSSM symmetry. The main requirement of this new minimum is zero cosmological constant which can be achieved if the fields a, p, ω, $\sigma(\bar{\sigma})$ take values such that the scalar potential $V = |W_{\phi_i}|^2/\Gamma'^2 = 0$ (where $\Gamma' = T + T^* - \frac{1}{3}(|p|^2 + 3|a|^2 + 6|\omega|^2 + |\sigma|^2 + |\bar{\sigma}|^2)$). The condition $W_{\phi_i} = 0$ required to have zero cosmological constant with broken SUSY (from the *vev* of the moduli fields T and T^*) in no-scale SUGRA is algebraically same as the condition for unbroken global supersymmetry in SUSY-SO(10) [24]. The field values a, p, ω, $\sigma(\bar{\sigma})$ which give $W_{\phi_i} = 0$ in SUSY SO(10) have been worked out in [24] and are given by

$$a = \frac{m}{\lambda} \frac{x^2 + 2x - 1}{1 - x}; \quad p = \frac{m}{\lambda} \frac{x(5x^2 - 1)}{(1 - x)^2};$$

$$\sigma \bar{\sigma} = \frac{2m^2}{\eta\lambda} \frac{x(1 - 3x)(1 + x^2)}{\eta(1 - x)^2}; \quad \omega = -\frac{m}{\lambda} x \quad (29)$$

where x is the solution of following cubic equation,

$$8x^3 - 15x^2 + 14x - 3 = -\frac{\lambda m_\Sigma}{\eta m}(1 - x)^2. \quad (30)$$

The soft SUSY breaking masses are proportional to the gravitino mass, which in no-scale SUGRA models with $V = 0$ is given [34,35] by

$$m_{3/2}^2 = e^G = e^K |W|^2. \quad (31)$$

In our case visible sector also contributes to gravitino mass as all the *vevs* are in units of m/λ so they can be of $O(M_P)$ from the inflationary conditions. However visible sector contribution can be made zero or negligible with field values of a, p, ω, $\sigma(\bar{\sigma})$ given by Eq. (29) and tuning $|W| \approx 0$. In that case only hidden sector and moduli fields determine the gravitino mass.

Also we need a pair of light Higgs doublets in MSSM. In the present scenario we have a 4×4 mass matrix $\mathcal{H}$ of MSSM Higgs doublets [36]. The form of mass matrix remains same as given in [36] with an extra factor of $1/\Gamma'$,

$$\mathcal{H} = \frac{1}{\Gamma'}$$

$$\times \begin{pmatrix} -m_H & \bar{\gamma}\sqrt{3}(\omega - a) & -\gamma\sqrt{3}(\omega + a) & -\bar{\gamma}\bar{\sigma} \\ -\bar{\gamma}\sqrt{3}(\omega + a) & 0 & -(2m_\Sigma + 4\eta(a + \omega)) & 0 \\ \gamma\sqrt{3}(\omega - a) & -(2m_\Sigma + 4\eta(a - \omega)) & 0 & -2\eta\bar{\sigma}\sqrt{3} \\ -\sigma\gamma & -2\eta\sigma\sqrt{3} & 0 & -2m + 6\lambda(\omega - a) \end{pmatrix}. \quad (32)$$

One out of the four Higgs doublets can be made light with the fine tuning condition of $Det\,\mathcal{H} = 0$. For fixed values of p, a, w, m, λ, it can be solved for m_H in terms of other free parameters of superpotential. For fixed real value of $x = -0.3471$ from $|W| \approx 0$ in the cases of successful inflation, m_H is given by

$$m_H = \frac{-0.887 \bar{\gamma}\gamma}{\eta} \text{ (case I)}; \quad m_H = \frac{-1.458 \bar{\gamma}\gamma}{\eta} \text{ (case IV)}. \quad (33)$$

For this m_H, one eigenvalue can be made light and the eigenvectors (left and right) corresponding to that eigenvalue can act as MSSM Higgs doublets.

3. Conclusions

In this work we show that the Starobinsky model of inflation can be derived from no-scale SUGRA SO(10) GUT for the specific intermediate symmetries of $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, $SU(5) \times U(1)$ and flipped $SU(5) \times U(1)$ gauge groups. The other intermediate symmetries $SU(4)_C \times SU(2)_L \times SU(2)_R$ or $SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B-L}$ do not give the slow-roll potential required for inflation. In the course of symmetry breaking topological defects like monopoles and cosmic strings can form. The defects formed in the first stage of symmetry breaking SO(10) $\rightarrow$ intermediate scale takes place during inflation and will be diluted away. After reheating when intermediate symmetry breaks to MSSM topological defects may form once again. The flipped $SU(5) \times U(1)$ and $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ breaking down to MSSM produces the cosmic strings [37] type of defect which is acceptable. However $SU(5) \times U(1)$ gives rise to monopoles after inflation and this case therefore can be ruled out from the consideration of topological defects in the cosmological evolution. The parameters of the SO(10) invariant superpotential are restricted by the requirement that the Starobinsky potential is obtained. These relations at the GUT scale can have testable consequences in the particle spectrum at low energy.

Acknowledgements

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