No scale SUGRA SO(10) motivated Starobinsky Model of Inflation

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We show that a Susy renormalizable theory based on gauge group SO(10) and Higgs system $10 \oplus 210 \oplus 126 \oplus \overline{126}$ 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 $SU(5) \times U(1)$ symmetry. The reheating in such scenario can occur via non perturbative decay of inflaton i.e. through “preheating”. After the end of inflation and preheating the universe finds a minimum which corresponds to MSSM, where Susy can be restored again.

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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 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 \( R^2 \) 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 was shown that the Starobinsky 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 inflation potentials which are identical to the Starobinsky potential at large field values. The natural choice for the inflaton in supergravity models are 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 \( \overline{5} \) Higgs in the superpotential has been constructed [18]. The SU(5) symmetry breaks to MSSM with the appropriate choice of vevs 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 minimal Kähler potential which gives polynomial potentials of inflation. However we show that a renormalizable Wess-Zumino superpotential of SO(10) GUT along with no-scale Kähler potential can give us Starobinsky kind inflation potential. The Higgs supermultiplets we consider are 10, 210, 126 (126). Among these the 210 and 126 (126) are responsible for breaking of SO(10) symmetry down to MSSM. 210 alone can give different intermediate symmetries [24] depending upon which of its MSSM singlet field is given a vev. Then 126 (126) breaks this intermediate symmetry to MSSM. The successful inflation potential can be achieved in the cases when the symmetry breaking pattern of the Higgs vevs 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 \( SU(5) \times U(1) \) symmetry. The other possible intermediate symmetries namely the Pati-Salam \( SU(4)_C \times SU(2)_L \times SU(2)_R \) or the \( SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B–L} \) gauge groups do not give phenomenologically correct inflation potentials.
After the end of inflation, the reheating can occur via non-perturbative decay of inflaton to bosons and fermions of the intermediate scale model. After reheating the system of universe finds a minimum which corresponds to MSSM where supersymmetry can be restored.

II. INFLATION IN SO(10) WITH NO SCALE SUGRA

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

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

The No-scale form of Kähler potential we take is

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

The $10$ and $\overline{126}$ are required for Yukawa terms to give masses to the fermions while $126(\overline{126})$ breaks the SO(10) gauge symmetry to MSSM together with $210$-plet. However to have an intermediate symmetry rather than MSSM, $210$-plet is sufficient. It can lead to various possible intermediate symmetries depending on which components of the $210$-plet get 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 as [29]

$$\Phi = 210 = (15, 1, 1) + (1, 1, 1) + (15, 1, 3) + (15, 3, 1) + (6, 2, 2) + (10, 2, 2) + (\overline{10}, 2, 2)$$

$$\Sigma = 126 = (\overline{10}, 1, 3) + (10, 3, 1) + (6, 1, 1) + (15, 2, 2)$$

$$\overline{\Sigma} = 126 = (\overline{10}, 3, 1) + (10, 1, 3) + (6, 1, 1) + (15, 2, 2)$$ (3)

The fields which will not break the MSSM symmetry are allowed to take vev. In this case they are

$$p = <\Phi(1,1,1)>, \quad a = <\Phi(15,1,1)>, \quad \omega = <\Phi(15,1,3)>, \quad \sigma = <\Sigma(10,3,1)>, \quad \bar{\sigma} = <\overline{\Sigma}(10,3,1)>$$ (4)

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

$$SO(10) \xrightarrow{210} \text{Intermediate symmetry} \xrightarrow{126} MSSM$$
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 $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ symmetry.

2. If $p \neq 0$ and $a=\omega=0$, this results to $SU(4)_C \times SU(2)_L \times SU(2)_R$ 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 $210$ is given as

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

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

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

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

The F-term potential has the following form,

$$V = e^G \left[ \frac{\partial G}{\partial \phi_i} K_{j*}^{ij} \frac{\partial G}{\partial \phi_j} - 3 \right]$$

Where

$$G = K + \ln W + \ln W^*$$

The kinetic term is given as $K_{j*}^{ij} \partial \phi^i \partial \phi_j$. Here $i$ runs over different fields $T,p,a$ and $\omega$. $K_{j*}^{ij}$ is the inverse of Kähler metric $K_{j*}^{ij}$ which is given as

$$K_{j*}^{ij} = \frac{1}{\Gamma^2} \begin{pmatrix} 3 & -p^* & -3a^* & -6\omega^* \\ -p & T + T^* - \frac{1}{3}(3|a|^2 + 6|\omega|^2) & a^*p & 2\omega^*p \\ -3a & ap^* & 3(T + T^*) - (|p|^2 + 6|\omega|^2) & 6a\omega^* \\ -6\omega & 2\omega p^* & 6a^*\omega & 6(T + T^*) - 2(|p|^2 + 3|a|^2) \end{pmatrix}$$

Where $\Gamma = T + T^* - \frac{1}{3}(|p|^2 + 3|a|^2 + 6|\omega|^2)$. After simplifying the potential (7) has the following form,

$$V = \frac{1}{\Gamma^2} \left| \frac{\partial W}{\partial \phi_i} \right|^2$$
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 $\chi_1$ direction. In Fig.1a we show $V(\chi_1, \chi_2=0, \chi_3)$ and in Fig.1b $V(\chi_1, \chi_2, \chi_3=0)$. We see that potential is flat along $\chi_1$ and confining along $\chi_2$ and $\chi_3$ respectively.

We assume that the non-perturbative Planck scale dynamics [10, 18, 30] fixes the value of $T = T^* = \frac{1}{2}$. After fixing the vev for $T$ the kinetic terms of $T$ and its mixed terms with other fields can be neglected. We then study all the possible cases of symmetry breaking mentioned earlier to study inflation 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 energy term is given as

$$L_{K.E.} = \frac{1}{(1-a^2)^2}((1-a^2)(\partial_\mu p)^2 + 3(\partial_\mu a)^2 + 6(1-a^2)(\partial_\mu \omega)^2)$$  \hspace{1cm} (11)

To get the canonical K.E. terms we need to do the transformation to new fields $\chi_1, \chi_2, \chi_3$.

$$a = \tanh\left(\frac{\chi_1}{\sqrt{3}}\right), \quad p = \text{sech}\left(\frac{\chi_1}{\sqrt{3}}\right)|\chi_2|, \quad \omega = \frac{1}{\sqrt{6}} \text{sech}\left(\frac{\chi_1}{\sqrt{3}}\right)|\chi_3|$$  \hspace{1cm} (12)

The potential $V(\chi_1, \chi_2, \chi_3)$ is flat along $\chi_1$ direction for $\chi_2=\chi_3=0$ and is confining in the orthogonal ($\chi_2, \chi_3$) directions as shown in Fig.1. So the potential in the limit $\chi_1 \neq 0$, $\chi_2 = \chi_3=0$ is

$$V = \frac{36m^2 \tanh^4 \left(\frac{\chi_1}{\sqrt{3}}\right) + 72m\lambda \tanh^3 \left(\frac{\chi_1}{\sqrt{3}}\right) + 36\lambda^2 \tanh^2 \left(\frac{\chi_1}{\sqrt{3}}\right)}{\left(1 - \tanh^2 \left(\frac{\chi_1}{\sqrt{3}}\right)\right)^2}$$  \hspace{1cm} (13)

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

$$V = 36m^2 (1 - e^{-2\chi_1/\sqrt{3}})^2$$  \hspace{1cm} (14)
The slow roll parameters for this potential are given as
\[ \eta = -\frac{8e^{-2\chi_1} \left( 1 - 2e^{-\frac{2\chi_1}{\sqrt{3}}} \right)}{3 \left( 1 - e^{-\frac{2\chi_1}{\sqrt{3}}} \right)^2}; \quad \epsilon = \frac{8e^{-4\chi_1}}{3 \left( 1 - e^{-\frac{2\chi_1}{\sqrt{3}}} \right)^2} \] (15)

Inflation ends when \( \eta \approx 1 \), which corresponds to field value of \( \chi_1 \approx 0.5 \). To have sufficient inflation which corresponds to \( N_{e-folds}=55 \) gives the field value \( \chi_1 \approx 4.35 \). The power spectrum for scaler perturbation \( P_R \) is given as
\[ P_R = \frac{V}{24\pi^2\epsilon} = \frac{9m^2\sinh^4 \left( \frac{\chi_1}{\sqrt{3}} \right)}{\pi^2} \] (16)

The value of \( P_R = (1.610\pm0.01) \times 10^{-9} \) given by Planck [7] requires value of \( m = 1.311 \times 10^{-6} \) in Planck units. The spectral index \( n_s = 0.964 \) and tensor to scalar perturbation ratio \( r = 0.002 \) for \( N_{e-fold} = 55 \).

• Case II: \( p \neq 0 \) and \( a=\omega=0, \ SU(4)_C \times SU(2)_L \times SU(2)_R \) symmetry.
In this case the old to new field transformation is given as,
\[ p = \sqrt{3}\tanh \left[ \frac{\chi_1}{\sqrt{3}} \right], \quad a = \frac{1}{\sqrt{3}} \sech \left[ \frac{\chi_1}{\sqrt{3}} \right]\chi_2, \quad \omega = \frac{1}{\sqrt{6}} \sech \left[ \frac{\chi_1}{\sqrt{3}} \right]\chi_3 \] (17)

Then the potential in the limit \( \chi_1 \neq 0, \chi_2 = \chi_3 = 0 \) is
\[ V = 3m^2 \sinh^4 \left( \frac{\chi_1}{\sqrt{3}} \right) \] (18)

This type of potential increases exponentially with \( \chi_1 \) and too steep to obey the slow roll conditions. The value of spectral index \( n_s \) has negative value over a wide range of field value, so 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.
In this case the old to new field transformation is given as,
\[ \omega = \frac{1}{\sqrt{2}} \tanh \left[ \frac{\chi_1}{\sqrt{3}} \right], \quad p = \sech \left[ \frac{\chi_1}{\sqrt{3}} \right]\chi_2, \quad a = \frac{1}{\sqrt{3}} \sech \left[ \frac{\chi_1}{\sqrt{3}} \right]\chi_3 \] (19)

Then the potential in the limit \( \chi_1 \neq 0, \chi_2 = \chi_3 = 0 \) is
\[ V = 72m^2 \sinh^2 \left[ \frac{\chi_1}{\sqrt{3}} \right] \cosh^2 \left[ \frac{\chi_1}{\sqrt{3}} \right] + \alpha \sinh^2 \left[ \frac{\chi_1}{\sqrt{3}} \right] \] (20)

Here \( \alpha = 5\lambda^2/8m^2 \). In this case for \( \alpha \geq -1 \) potential increases exponentially with \( \chi_1 \) so give similar results to case II. For \( \alpha <-1 \) potential energy becomes negative for \( \chi_1 \gtrsim 1 \) and grows with large values of \( \chi_1 \).
• 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 is given as

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

(21)

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

(22)

After making the K.E. term canonical with transformation \( x = \sqrt{\frac{2}{10}} \tanh \frac{\chi}{\sqrt{3}} \), the potential we get is

\[
V = 55.2m^2(1 - e^{-\frac{2\chi}{\sqrt{3}}})^2
\]

(23)

for \( \lambda = -\frac{1}{3}\sqrt{\frac{10}{3}}m \), which is Starobinsky kind of potential for inflation with different relation among superpotential parameters \( m \) and \( \lambda \) than the case I. In this case value of \( m = 1.06 \times 10^{-6} \) is required to satisfy the constraints from CMB observations mentioned in case I.

At the end of inflation the preheating [31] can occur via non perturbative decay of inflaton \( \chi_1 \) to scalar bosons which have a trilinear term with \( \Phi \) in superpotential e.g. \( \Phi H (\gamma \Sigma + \bar{\gamma} \bar{\Sigma}) \). Then the \( K_{\Sigma}^\nu |W_\Sigma|^2 \) and \( K_{\bar{\Sigma}}^\nu |W_{\bar{\Sigma}}|^2 \) type of terms gives us

\[
V \supset (|\gamma|^2 |\bar{\gamma}|^2 |H|^2 + |\gamma|^2 |\Sigma|^2 + |\bar{\gamma}|^2 |\bar{\Sigma}|^2)|\sinh\left(\frac{\chi_1}{\sqrt{3}}\right)|^2
\]

(24)

Near the origin \( \sinh\left(\frac{\chi_1}{\sqrt{3}}\right) \approx \frac{\chi_1}{\sqrt{3}} \),

\[
V \supset (|\gamma|^2 |\bar{\gamma}|^2 |H|^2 + |\gamma|^2 |\Sigma|^2 + |\bar{\gamma}|^2 |\bar{\Sigma}|^2)\left(\frac{\chi_1}{\sqrt{3}}\right)^2
\]

(25)

The 210 inflaton produces the 10 and \( 126 \) which have Yukawa couplings with the fermions and will decay into the SM fermions and the right-handed neutrino to give a radiation dominated universe at the end of inflation. \( \gamma, \bar{\gamma} \) are free parameters of the superpotential, so they can be set to have effective preheating with reheat temperature \( T_R \sim 10^{13} GeV \) needed for leptogenesis [32].

During inflation Susy is broken at Planck scale and we need to restore Susy at the TeV scales to explain the light Higgs of the standard model. This can be done if the universe finds a minimum where the field values are given by [24]

\[
a = \frac{m x^2 + 2x - 1}{\lambda \left(1 - x\right)} \quad ; \quad p = \frac{m x(5x^2 - 1)}{\lambda \left(1 - x\right)^2} \quad ; \quad \sigma \bar{\sigma} = \frac{2m^2 x(1 - 3x)(1 + x^2)}{\eta \lambda \left(1 - x\right)^2} \quad ; \quad \omega = -\frac{m}{\lambda} x
\]

(26)
Where \( x \) is the solution of following cubic equation

\[
8x^3 - 15x^2 + 14x - 3 = \frac{\lambda m_\Sigma}{\eta m}(1 - x)^2
\]  

(27)

In case of SUGRA an additional condition of \( W=0 \) is required to have MSSM symmetry. Putting these field values in the expression of superpotential and setting \( W = 0 \) gives us a fix value of \( x \). Since all the vevs are in units of \( m/\lambda \) so they can be of \( O(M_{pl}) \) from inflation conditions. However the main requirement of MSSM is a pair of light Higgs. In the present scenario we have a \( 4 \times 4 \) mass matrix \( \mathcal{H} \) of MSSM Higgs doublets [33]. However after making the kinetic terms canonical the fields in the mass matrix will be in new basis say \( \chi_1, \chi_2 \) etc., but the form of mass matrix remains similar to given in [33],

\[
\mathcal{H} = \begin{pmatrix}
-m_H & \gamma \sqrt{3}(w - a) & -\gamma \sqrt{3}(w + a) & -\gamma \bar{\sigma} \\
-\gamma \sqrt{3}(w + a) & 0 & -(2m_\Sigma + 4\eta (a + \omega)) & 0 \\
\gamma \sqrt{3}(w - a) & -(2m_\Sigma + 4\eta (a - \omega)) & 0 & -2\eta \bar{\sigma} \sqrt{3} \\
-\sigma \gamma & -2\eta \sigma \sqrt{3} & 0 & -2m + 6\lambda (\omega - a)
\end{pmatrix}
\]  

(28)

One out of the four Higgs doublets can be made light with the fine tuning condition of \( \text{Det}\mathcal{H}=0 \). For fixed values of \( p, a, w, m, \lambda \) it can be solved for \( m_H \) in terms of other free parameters of superpotential. For real value of \( x=-0.3471 \) in all cases, \( m_H \) is given as,

\[
m_H = \frac{-0.887\gamma \gamma}{\eta} \quad \text{(case I)}; \quad m_H = \frac{-1.458\gamma \gamma}{\eta} \quad \text{(case IV, V)}
\]  

(29)

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

### III. CONCLUSIONS

In this work we show that the Starobinsky model of inflation can be derived from no-scale SUGRA SO(10) GUT for specific intermediate symmetries namely the \( SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \) and the \( SU(5) \times U(1) \) and flipped \( SU(5) \times U(1) \) gauge groups. The other intermediate symmetries namely the \( SU(4)_C \times SU(2)_L \times SU(2)_R \) or the \( SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B-L} \) gauge groups do not give the slow-roll potentials required for inflation. 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.

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