Inflation With A Realistic SO(10) Model

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Abstract. We implement inflation within a realistic supersymmetric SO(10) model in which the doublet-triplet splitting is realized through the Dimopoulos-Wilczek mechanism, the MSSM $\mu$ problem is resolved, and higgsino mediated dimension five nucleon decay is heavily suppressed. The cosmologically unwanted topological defects are inflated away, and from $\delta T/T$, the $B - L$ breaking scale is estimated to be of order $10^{16} - 10^{17}$ GeV. Including supergravity corrections, the scalar spectral index $n_s = 0.99 \pm 0.01$, with $|dn_s/d\ln k| \lesssim 10^{-3}$.

Keywords: Inflation, $\delta T/T$, SO(10)

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In a class of supersymmetric (SUSY) models, inflation is associated with spontaneous breaking of a gauge symmetry, such that $\delta T/T$ is proportional to $(M/M_{\text{Planck}})^2$, where $M$ denotes the symmetry breaking scale and $M_{\text{Planck}} (\approx 1.2 \times 10^{19}$ GeV) denotes the Planck mass [2, 3]. Thus, from measurements of $\delta T/T$, $M$ is estimated to be of order $10^{16}$ GeV [2, 4]. The scalar spectral index $n_s$ in these models is very close to unity in excellent agreement with recent fits to the data [5]. A $U(1)$ $R$-symmetry plays an essential role in the construction of these inflationary models. These models possess another important property, namely with the minimal Kähler potential, the supergravity (SUGRA) corrections do not spoil the inflationary scenario [3], which has been realized with a variety of attractive gauge groups including $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ ($\equiv G_{LR}$) [6], $SU(4)_c \times SU(2)_L \times SU(2)_R$ ($\equiv G_{422}$) [7], $SU(5) \times U(1)$ [8], and $SU(5)$ [9]. We aim to implement inflation within a realistic SO(10) model [1].

SO(10) has two attractive features, namely, it predicts the existence of right handed neutrinos as well as the seesaw mechanism. These two features are very helpful in understanding neutrino oscillations and also in generating a baryon asymmetry via leptogenesis. Furthermore, it seems easier to realize doublet-triplet (DT) splitting without fine tuning in SO(10) (say via the Dimopoulos-Wilczek mechanism [10]) than in SU(5).

To implement SO(10) inflation we would like to work with a realistic model with the following properties: DT splitting is realized without fine tuning, and the low energy theory coincides with the minimal supersymmetric standard model (MSSM). The MSSM $\mu$ problem should also be resolved, and higgsino mediated dimension five ($d = 5$) nucleon decay should be adequately suppressed. Gauge boson mediated nucleon decay is still present with a predicted nucleon lifetime of order $10^{34} - 10^{36}$ yrs. Finally, matter parity is unbroken, so that the LSP is stable and makes up the dark matter in the universe. To achieve natural DT splitting and the MSSM at low energies with SO(10), we will follow Refs. [11, 12], with suitable modifications needed to make the scheme consistent with the desired inflationary scenario, and also to avoid potential cosmological problems (monopoles, moduli, etc.). While doing this we would like to also ensure that the SUGRA corrections also do not disrupt the inflationary scenario.

A minimal set of Higgs required to break SO(10) to the MSSM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ ($\equiv G_{SM}$) is $45_H$, $16_H$, $\overline{10}_H$. A non-zero vacuum expectation value (VEV) of $45_H$ along the $B - L$ ($h_{LR}$) direction breaks SO(10) to $G_{LR}$ ($SU(4)_c \times SU(2)_L \times U(1)_R$) and produces magnetic monopoles. The $16_H$, $\overline{10}_H$ VEVs break SO(10) to SU(5) and induce masses for the right handed neutrinos via $d = 5$ operators. One of our goals is to make sure that the topological defects do not pose cosmological difficulties. Thus, it would be helpful if during inflation SO(10) is, for instance, broken to $G_{LR}$, $SU(4)_c \times SU(2)_L \times U(1)_R$, or $G_{SM}$.

To implement DT splitting without fine tuning and eliminate $d = 5$ proton decay, and to recover the MSSM at low energies with the $\mu$ problem resolved, we need an additional $45$-plet ($45_H$), two additional $16 + \overline{10}$ pairs, two $10$-plets ($10_H$ and $10$), and several singlets [11, 12]. One more $45$-plet is also required by $U(1)_{R}$-symmetry. This symmetry, among other things, plays an essential role in realizing inflation, and its $Z_2$ subgroup coincides with the MSSM matter parity. The SO(10) singlet superfields are denoted as $S$, $X$, $X'$, $Y$, $P$, $\mathcal{P}$, $Q$, and $\mathcal{Q}$, whose

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roles will be described below. Table 1 displays the quantum numbers (under the global $U(1)_R$ and $U(1)_A$ symmetries) of all the Higgs sector superfields and the third family matter field ($16_Y$). We will take the same notation for the superfields and their scalar components.

To break $SO(10)$ to $G_{LR}$, consider the superpotential,

\[
W_{45} = \frac{\alpha}{6M_e} X' Y \Tr (45^2) - \frac{\beta}{6} Y \Tr (4545^H) + \frac{\gamma}{56M_e} Y \Tr (4545^H) + \frac{\kappa}{6M_e} \Tr (4545^3) H,
\]

where $\alpha, \beta, \gamma, \kappa$ are dimensionless parameters, and $M_e$ ($\sim 10^{16}$ GeV) denotes the cutoff scale. As will be explained, $X, X'$, and $Y$ can develop non-zero VEVs, $\langle X' \rangle \sim \langle Y \rangle \sim 10^{16}$ GeV. Due to non-zero $\langle Y \rangle$, $45^H$ can also obtain a VEV in the $B-L$ direction from the $\beta$ and $\gamma$ terms of Eq. (1).

\[
\langle 45^H \rangle = \text{diag} (v, v; 0, 0) \otimes i \sigma_2
\]

From the $\kappa$ and $\rho$ terms, and the “D-term” potential, $16_H$ and $\overline{16_H}$ develop VEVs of order $M_{B-L}$, breaking $SO(10)$ to $SU(5)$,

\[
|\langle 16_H \rangle|^2 = |\langle \overline{16_H} \rangle|^2 = \frac{M_{B-L}^2}{2\zeta} \left[ 1 + \sqrt{1 - 4\zeta} \right],
\]

where $\zeta \equiv \rho M_{B-L}/(\kappa M^2_e)$ [7], while $\langle S \rangle \approx 0$ up to corrections of $O(m_{3/2}^2)$ by including soft SUSY breaking terms in the scalar potential [6]. Together with Eq. (2), the $SO(10)$ is broken to the $G_{SM}$. The MSSM Higgs doublets arise from $10_H$. With $\langle S \rangle \approx -m_{3/2}/\kappa$, the $\mu$ term from Eq. (3) is of order $\langle \lambda / \kappa \rangle m_{3/2} \sim \text{TeV}$.

When $SO(10)$ breaks to $G_{SM}$ by an adjoint and a vector-like pair of spinorial Higgs, the superfields associated with $\langle (3, 2)_1 \rangle, \langle (3, 1)_1 \rangle$, and $\langle \overline{1, \overline{3}}_1 \rangle$ h.c. from the $45^H$ and $16_H + \overline{16_H}$ turn out to be pseudogoldstone modes [11]. Such extra light multiplets would spoil the unification of the MSSM gauge couplings, and therefore must be eliminated. The simplest way to remove them from the low energy spectrum is to introduce couplings such as $16_H 45^H \overline{16_H}$. However, it destabilizes the form of $45^H_H$ given in Eq. (2), in such a way that at the SUSY minimum, $v = 0$ is required. It was shown in Ref. [11] that with the ‘$\lambda$’ couplings ($i = 1, 2, 3, 4$) and an additional $16 - \overline{16}$ pair in Eq. (3), the unwanted pseudogoldstone modes all become superheavy, keeping intact the form of Eq. (2) at the SUSY minimum.

From the “F-flat conditions” with $16_H$ and $\overline{16_H}$ acquiring non-zero VEVs, one finds

\[
\langle 45^H \rangle / \langle Y \rangle = \frac{\lambda_2}{\lambda_1} (P^2), \quad \langle 45^H \rangle / \langle Q \rangle = \frac{\lambda_4}{\lambda_3} \Tr (45^H)^2.
\]

Thus, if $P$ and $Q$ develop VEVs, $\langle 45^H \rangle$, $\langle 45_{H}^F \rangle$, and $\langle Y \rangle$ should also appear. We will soon explain how $\langle P \rangle$ and $\langle Q \rangle$ arise. Since $\langle Y \rangle$ is related to $\langle 45^H \rangle$ via Eq. (2), both are uniquely determined. We assume that $\langle 45_H^H \rangle$ points in

\[
\begin{array}{|c|c|c|c|c|c|c|c|c|}
\hline
& S & X & X' & Y & P & \overline{P} & Q & \overline{Q} \\
\hline
R & 1 & -1 & -1 & 0 & 0 & 0 & -2 & 2 \\
A & 0 & -2/3 & -2/3 & -1/3 & -1/4 & 1/4 & -1/2 & 1/2 \\
\hline
& 16 & 16' & 16' & 16_H & \overline{16_H} & 16_H & 45 & 45_H & 45_H \\
\hline
R & 1 & 3 & 2 & 2 & 0 & 0 & 1/2 & 1 & 0 & -1 \\
A & 1/2 & 2/3 & 2/3 & 2/3 & 0 & 0 & 1/2 & -1/6 & -1/3 \\
\hline
\end{array}
\]

\[\begin{array}{|c|c|c|c|c|c|c|c|c|}
\hline
16 & 16' & 16' & 16_H & \overline{16_H} & 16_H & 45 & 45_H & 45_H \\
\hline
1 & 3 & 2 & 2 & 0 & 0 & 1/2 & 1 & 0 & -1 \\
1/2 & 2/3 & 2/3 & 2/3 & 0 & 0 & 1/2 & -1/6 & -1/3 \\
\end{array}\]

\[\begin{array}{|c|c|c|c|c|c|c|c|c|}
\hline
\end{array}\]

\[\begin{array}{|c|c|c|c|c|c|c|c|c|}
\hline
\end{array}\]
Recall that \((45'_H)\) is employed to suppress higgsino mediated \(d = 5\) nucleon decay [12]. Similarly, due to the presence of the \(\lambda_k^i\) \((i = 5, 6, 7, 8)\) couplings in Eq. (3), the low energy spectrum is protected even with the \(45'_H\) present [12]. With non-zero VEVs for \(45'_H\) and \(Y, X\) and \(X'\) slide to the values satisfying

\[
\langle 45'_H \rangle = \text{diag} (0, 0, 0; v', v') \otimes i\sigma_2 .
\]  

(6)

In order to guarantee the \(\lambda_k^i\) couplings in Eq. (3) and to forbid \(16_H 45'_H 16_H\), the \(U(1)\) symmetries in TABLE 1 are essential.

To obtain non-vanishing VEVs for \(P\) and \(Q\), one could consider the following superpotential,

\[
W_{PQ} = S \left[ \kappa_1 P \bar{P} + 2 \kappa_2 Q \bar{Q} \right] - \frac{S}{M_5} \left[ \rho_1 (P \bar{P})^2 + \rho_2 (Q \bar{Q})^2 \right] ,
\]  

(8)

such that \(\langle P \bar{P} \rangle = (\kappa_1 / \rho_1) M_5^2 \sim \langle Q \bar{Q} \rangle = (\kappa_2 / \rho_2) M_5^2 \sim M_{GUT}^2\). The \(\lambda_{2,3}\) terms in Eq. (3) just determine \(45'_H\), \(\langle Y \rangle\), and \(\langle 45'_H \rangle\). With the inclusion of soft SUSY breaking terms, \(\langle P \rangle\), \(\langle \bar{P} \rangle\), \(\langle Q \rangle\), and \(\langle \bar{Q} \rangle\) would be completely fixed. To avoid potential cosmological problems associated with moduli fields, we assume that the VEVs satisfy the constraints \(\langle P \rangle = \langle \bar{P} \rangle\) and \(\langle Q \rangle = \langle \bar{Q} \rangle\). This could be made plausible by assuming universal soft scalar masses, and that the SUSY breaking “\(A\)-terms” asymmetric under \(P \leftrightarrow \bar{P}\) and \(Q \leftrightarrow \bar{Q}\) are small enough. Since the fields that couple to \(P, \bar{P}, Q\) and \(\bar{Q}\) are all superheavy, the soft parameters would be radiatively stable at low energies. Thus, at the minimum of the scalar potential, we have four mass eigen states, \(\langle P \pm \bar{P} \rangle/\sqrt{2} = (\pm P_+\) and \(\langle Q \pm \bar{Q} \rangle/\sqrt{2} = (\pm Q_+).\) While \(P_+\) and \(Q_+\) obtain superheavy masses and large VEVs of order \(M_{GUT}\), \(P_-\) and \(Q_-\) remain light \((\sim m_{3/2})\) with vanishing VEVs.

With \(45'_H\) in Eq. (2), the “\(D\)T splitting problem” can be resolved [10]. Consider the superpotential

\[
W_{10} = y_1 1045'_H 10 + y_2 1045'_H 10_h .
\]  

(9)

From the \(y_1\) term, only the doublets contained in \(10\) become superheavy [12], and from the \(y_2\) term only the color triplet fields included in \(10\) and \(10_h\) acquire super-heavy masses [10, 11, 12]. Since the two color triplets contained in \(10_h\) do not couple at all in Eq. (9), \(d = 5\) nucleon decay is eliminated in the SUSY’ limit [12]. Note that operators such as \(1010_h, 10^2, [1010^2]_H^T (45'_H)\) and so on are allowed by \(SO(10)\) and, unless forbidden, would destroy the gauge hierarchy. The \(U(1)\) symmetries in TABLE 1 are once again crucial in achieving this.

Although the superpotential coupling \((S) 10^2\) induces higgsino mediated \(d = 5\) nucleon decay, there is a suppression factor of \(m_{3/2} / M_{GUT}\). Thus, we expect that nucleon decay is dominated by the exchange of the superheavy gauge bosons with an estimated lifetime \(\tau_p \rightarrow e^+ e^-\) of order \(10^{34} \sim 10^{36}\) yrs. Note that we have assumed that \(d = 5\) operators such as \(16, 16, 16, 16, 16, 16, 16_H\) and so on, where the subscripts are family indices of the matter, are adequately suppressed by assigning suitable \(R\) and \(A\) charges to these matter fields.

Consider next the superpotential couplings involving the third generation matter superfields,

\[
W_m = y_3 16, 16, 16, 10_h + \frac{y_v}{M_5} 16, 16, 16_H, 16_H .
\]  

(10)

The first term yields Yukawa unification so that the MSSM parameter \(\tan \beta \approx m_h / m_{\tilde{t}_R}\). From the \(y_v\) term, the right handed neutrino masses are \(\lesssim y_v M_{Planck}^2 / M_5 \sim 10^{14}\) GeV. Right handed neutrino masses of order \(10^{15}\) GeV and smaller can yield a mass spectrum for the light neutrinos through the seesaw mechanism, that is suitable for neutrino oscillations. The role of ‘matter parity’ is played by the unbroken \(Z_2\) subgroup of the \(U(1)\) \(R\)-symmetry [6]. Thus the LSP in our model is stable and contributes to the dark matter in the universe.

Let us now discuss how inflation is implemented in the model described so far. The “\(F\)’term” scalar potential in SUGRA is given by

\[
V_F = e^{K / M_P^2} \left[ \sum_{i,j} (K^{-1})_{ij} (D_\partial W)(D_\partial W) - \frac{3 |W|^2}{M_P^2} \right] ,
\]  

(11)

where \(M_P = (M_{Planck}) / \sqrt{8\pi} = 2.4 \times 10^{18}\) GeV) denotes the reduced Planck mass. \(K\) and \(W\) are the Kähler potential and the superpotential, respectively. \((K^{-1})_{ij}\) in Eq. (11) denotes the inverse of \(\partial^2 K / \partial \partial \partial \phi^i \partial \phi^j\). In our case, \(K\) is composed of Eqs. (1), (3), (8), (9), and (10). \(D_\partial W\) is defined as \(\partial W / \partial \phi^i + (\partial K / \partial \phi^i)(W / M_P^2)\). The Kähler potential could be expanded as \(K = |\phi|^2 + c_4 |\phi|^4 / M_P^2 + \cdots\). Here, we consider the minimal case with \(\partial^2 K / \partial \partial \partial \phi^i \partial \phi^j = \delta^i_j\). Indeed, higher order terms in \(K\) (with \(c_4 \lesssim 10^{-2}\)) do not seriously affect inflation [3].
1/72 [7], and $\langle 16_H \rangle$, $\langle \overline{16}_H \rangle$, $\langle P \rangle$, $\langle P \rangle$, $\langle Q \rangle$, $\langle Q \rangle \neq 0$ with the inflationary superpotential given by [7],

$$W_{	ext{inf}} \approx -\kappa S \left[ M_{\overline{B}-L}^2 - 16_H \overline{16}_H + \frac{\rho}{\kappa M_{\overline{B}-L}^2} (16_H \overline{16}_H)^2 \right]$$

$$- \frac{k_1}{\kappa} \langle P \rangle^2 + \frac{\rho_1}{\kappa} \langle Q \rangle^2 - \frac{k_2}{\kappa} \langle Q \rangle \overline{Q} + \frac{\rho_2}{\kappa M_{\overline{B}-L}^2} \langle Q \rangle^2 \right]$$

$$\equiv -\kappa S M_{\text{eff}}^2,$$  \hspace{1cm} (12)

where $M_{\text{eff}}^2$ turns out to be of order $M_{\overline{B}-L}^2$. With $D_3 W \approx -\kappa M_{\text{eff}}^2(1 + |S|^2/M_{\overline{B}-L}^2)$, Eq. (11) becomes

$$V_F \approx \left( 1 + \sum_i |\langle \phi_i \rangle|_T^2 \right) \left[ \kappa^2 M_{\text{eff}}^2 \left( 1 + \frac{|S|^2}{2M_{\overline{B}-L}} \right) + \left( 1 + \frac{|S|^2}{2M_{\overline{B}-L}} \right) \sum |\langle \phi_i \rangle|_T^2 \right],$$  \hspace{1cm} (13)

where all scalar fields except $S$ contribute to $\phi_i$. In Eq. (13) the quadratic term of $S$ from $|D_3 W|^2$, which is of order $(\kappa M_{\text{eff}}^2/M_{\overline{B}-L})^2$, has canceled out with the factor $-3W_{\text{inf}}^2/M_{\overline{B}-L}^2$ and the quadratic term in $S$ from “$e^{K/M_{\overline{B}-L}^2}$” is a common feature in this class of models [3]. Thus, the dominant mass term for $S$ is

$$V_S \approx \sum_i |\langle \phi_i \rangle|_T^2 \sim \left( \frac{M_{\text{GUT}}}{M_{\overline{B}-L}} \right)^2 \times H^2 |S|^2,$$  \hspace{1cm} (14)

where $\phi_i = X(i), Y, 45_H, 45_H^*$, and $H (\approx M_{\text{eff}}^2/M_{\overline{B}-L})$ denotes the “Hubble induced mass.” Such a small mass term of $S \langle H^2 |S|^2 \rangle$ does not spoil the slow roll conditions. Note that the $U(1) R$-symmetry ensures the absence of $S^2$, $S^3$, etc., in the superpotential, which otherwise could spoil the slow-roll conditions.

At one of the local minima, $\langle 16_H \rangle$, $\langle \overline{16}_H \rangle$, $\langle P \rangle$, $\langle P \rangle$, $\langle Q \rangle$, and $\langle Q \rangle$ acquire the non-zero VEVs; $|\langle 16_H \rangle|^2 \approx |\langle \overline{16}_H \rangle|^2 \approx \kappa^2 M_{\overline{B}-L}^2 / (2p_1)$, $|\langle P \rangle|^2 \approx |\langle Q \rangle|^2 \approx \kappa^2 M_{\overline{B}-L}^2 / (2p_2)$, and $|\langle Q \rangle|^2 \approx |\langle Q \rangle|^2 \approx \kappa^2 M_{\overline{B}-L}^2 / (2p_2)$ [7]. Since $P$ and $Q$ develop VEVs, $\chi^0(i), Y, 45_H, 45_H^*$ should also achieve VEVs from $D_{16}W = D_{\overline{16}}W = 0$ even during inflation. Consequently, $SO(10)$ and $U(1)_A$ are broken to $G_{\text{SM}}$ during inflation. Note that $\langle P \rangle = \langle P \rangle$ and $\langle Q \rangle = \langle Q \rangle$ lead to $\langle P \rangle = \langle Q \rangle = 0$. Since $\langle P \rangle$ and $\langle Q \rangle$ vanish both during and after inflation, oscillations by such light ($< m_{3/2}$) scalars would not arise after inflation has ended. A non-zero vacuum energy from the “$e^{-} term potential induces universal “Hubble induced scalar mass terms” ($\approx \kappa^2 M_{\text{eff}}^2/M_{\overline{B}-L}^2 \times |\phi_i|^2$), which are read off from Eq. (13). But such small masses ($< m_{3/2}$) cannot not much affect the VEVs of the superheavy scalars of order $M_{\text{GUT}}$.

With SUSY broken during inflation ($f_S \neq 0$), there are radiative corrections from the $16_H, \overline{16}_H$ supermultiplets, which provide logarithmic corrections to the tree level potential $V_F \approx \kappa^2 M_{\text{eff}}^2 \approx \kappa^2 M_{\overline{B}-L}^2 (4\zeta)^2$, and thereby drive inflation [2]. In our model, the scalar spectral index turns out to be $n_s = 0.99 \pm 0.01$ for $\kappa < 10^{-2}$, and the symmetry breaking scale $M_{\overline{B}-L}$ is estimated to be around $10^{10} \sim 10^{17}$ GeV [4]. When inflation is over, the inflatons decay into right handed neutrinos. Following Ref. [13], the lower bound on $T_r$ is $T_r \approx 10^9$ GeV for $\kappa < 10^{-2}$. The inflaton decay into right handed neutrinos yields the observed baryon asymmetry via leptogenesis [14]. Assuming non-thermal leptogenesis and hierarchical right handed neutrinos, we estimate the three right handed neutrinos masses to be of order $10^{14}$ GeV, $(10-20) \times T_r$, and few $\times T_r$. Note that with $\kappa < 10^{-2}$ the inflaton (with mass $\approx \sqrt{\kappa^2 M_{\overline{B}-L}^2}$) can not decay into the heaviest right handed neutrino (of mass $\approx 10^{14}$ GeV). Thus, the latter does not play a direct role in leptogenesis.

In summary, our goal here was to realize inflation in a realistic SUSY $SO(10)$ model. A global $U(1)_A$ and the $U(1) R$-symmetry plays essential roles in the analysis. The scalar spectral index is $n_s = 0.99 \pm 0.01$, which will be tested by several ongoing experiments. Proton decay proceeds via $e^+e^0$, with an estimated lifetime of order $10^{24} \sim 10^{36}$ yrs. The LSP is stable. While the heaviest right-handed neutrino weighs around $10^{14}$ GeV, the one primarily responsible for non-thermal leptogenesis has mass of order $10 T_r$, where the reheat temperature $T_r$ is around $10^9 \sim 10^9$ GeV.

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