Non-Singlet Inflaton in Supergravity

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Abstract

We consider the case where the inflaton is non-singlet in a supergravity framework. The η-problem is avoided by defining a shift symmetry on the charged inflaton fields in a consistent way with the gauge symmetry. We review two scenarios, one of them depends on a $U(1)$ gauge symmetry group and the other depends on flipped GUT gauge symmetry.

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1. INTRODUCTION

Inflation paradigm is the most acceptable framework that provides a solution to the problems associated with the original Big Bang model, such as flatness, horizon, and monopole problems. On the other hand supersymmetry is a good candidate for physics beyond the standard model of elementary particle physics that has good candidates for the inflaton fields. In such a high scale of inflation physics, supergravity corrections should be taken in to account.

However in a supergravity framework, a problem arises, called the $\eta-$problem, due to contributions to the inflaton mass of order of Hubble scale. Defining a shift symmetry can overcome such a problem in supergravity. On the other hand, requiring that the inflaton is charged under a gauge group ivites us to define the shift symmetry in a consistent way as will be advocated in two prominent examples below.

The paper is organized as follows. In section 2 we present a model with a $U(1)$ charged inflaton and describe the inflation observables. In section 3 we show that a charged inflaton can be embedded in a flipped GUT (FGUT) scenario, then display the observables in different cases of the waterfall regime. Finally we discuss the reheating in both cases.

2. $U(1)$ CHARGED INFLATON IN SUPERGAVITY

In this section we discuss the possibility if realizing the infation scenario in the direction of a pair of superfields that are $U(1)$ oppositely charged. Let's begin with the most general superpotential that is renormalizable and consistent with R-symmetry, containing two superfields ϕ_1 and ϕ_2 carrying opposite charges under a $U(1)$ gauge symmetry, and a singlet S [2]

$$W = \lambda S(\phi_1\phi_2 + M^2), \tag{1}$$

In that case one should impose a shift symmetry [7, 8, 9] on the fields ϕ_1 and ϕ_2, in a consistent way to avoid the troublesome η-problem appearing in supergravity models of inflation. Hence, the Kähler potential will have the form [2]

$$K = |\phi_1 + \overline{\phi_2}|^2 + |S|^2 - \zeta|S|^4, \tag{2}$$

Here, ζ plays important role in determining the inflaton multiplet, S or $\phi_{1,2}$. If ζ is negative, one has a scenario similar to [5] where the supergravity corrections are taken into account in the hybrid inflation scenario of Ref. [4]. In that case, the inflaton comes from the S multiplet while the fields $\phi_{1,2}$ will play the role of the waterfall fields and will be frozen at the origin during the inflation. It is worth mentioning that the R-symmetry protects the inflaton mass against large supergravity corrections in such scenario. Hence, the η-problem is absent. This case is a small-field inflation and here M is of order GUT scale.

On the other hand, considering ζ to take positive values is more motivated if one interprets the quartic term of S in the Kähler as loop corrections [6]. In that scenario, the inflaton comes from the $\phi_{1,2}$ whereas S represents the stabilizer field [10, 11].

2.1. *Non-anomalous $U(1)$ case*

In a non-anomalous $U(1)$ symmetry, the F-term and D-term scalar potentials are given by

$$V_F = e^K \left[K^{I\bar{J}} D_I W \overline{D_J W} - 3|W|^2\right], \tag{3}$$

$$V_D = \frac{g^2}{2}\left(|\phi_1|^2 - |\phi_2|^2\right)^2, \tag{4}$$

where I, J run over the superfields $\{S, \phi_1, \phi_2\}$ and $D_I W = \partial_I W + W \partial_I K$. Here, we work in units where the reduced Planck mass $M_p = 1$. Using the following parametrizations for the complex fields,

$$\begin{aligned} S &\equiv s + i\sigma, \\ \phi_1 + \bar{\phi}_2 &\equiv \alpha + i\beta \\ \phi_1 - \bar{\phi}_2 &\equiv \rho\, e^{i\theta/2M}. \end{aligned} \tag{5}$$

Hence the SUSY minimum of the potential is given by

$$\langle s \rangle = \langle \sigma \rangle = \langle \alpha \rangle = \langle \beta \rangle = 0 \quad \text{and} \quad \langle \rho \rangle = 2M. \tag{6}$$

It is clear that the Kähler potential (2) doesn't depend on the fields ρ and θ. However, θ is massless in the vacuum and is corresponding to unphysical degree of freedom.[1] In that case the slow rolling inflaton will be the field ρ, whereas the other real components are fixed at zero during the inflation. Accordingly, the effective inflationary potential is given by

$$V_{\text{inf}}(\rho) = \frac{\lambda^2}{16}\left(\rho^2 - 4M^2\right)^2. \tag{7}$$

This potential has been studied in details in [13, 10, 12]. If M has values $M < 1$, then the potential is mainly governed by the quartic term which is excluded by WMAP and Planck observations. On The other hand, for $M \gg 1$, we have two possible regimes for inflation.

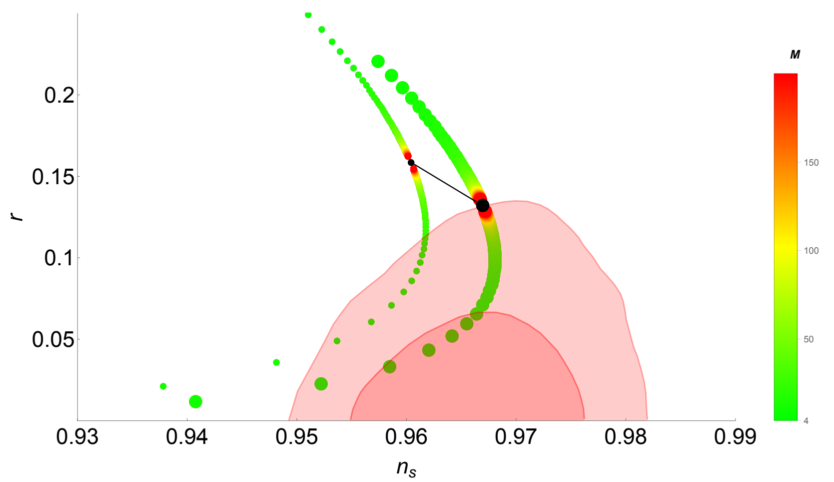

FIGURE 1: The $n_s - r$ observables of the inflationay potential Eq.(7), where we have sanned over the parameter M from $4 - 200M_p$. The large points correspond to 60 e-folds while the smaller points correspond to 50 e-folds. The black ones correspond to the observables of $m^2\phi^2/2$ chotic inflation. Light red regions correspond to the 1 and 2 sigma exclusion limits released by the Planck collaboration [1], (Planck TT + LowP).

Fig. 1 shows the inflationary observables: the spectral index n_s versus the tensor to scalar ratio r. The two black points correspond to the chaotic inflation $V = m^2\phi^2/2$, observables. The points above and below the black points correspond to two different inflationary regimes of the potential (7). Points below the black one correspond to the regime when the inflaton rolls from initial values less than $2M$. While the above ones, correspond to the regime when the inflaton rolls from initial values larger than $2M$.

2.2. *Anomalous $U(1)$ case*

Fayet Illiopoulos (FI) term can exist in case we have anomalous $U(1)$. The D-term potential will be modified as follows

$$V_D = \frac{g^2}{2}\left(|\phi_1|^2 - |\phi_2|^2 + \xi\right)^2, \tag{8}$$

and the SUSY minimum turns out to be

$$|\phi_{1,2}|^2 = \frac{\mp\xi + \sqrt{\xi^2 + 4M^4}}{2}, \tag{9}$$

with the fields s and σ remain stabilized at zero. Working in the basis

$$\phi_1 \pm \bar{\phi}_2 \equiv \alpha_\pm + i\beta_\pm, \tag{10}$$

[1] It is the Nambu Goldstone boson arising after spontaneous breaking of the $U(1)$ gauge symmetry.

we find that the inflaton will correspond to α_- and other spectator fields will be frozen during the inflation to [2]

$$\begin{aligned}
\langle s\rangle_{inf} &= 0\,,\\
\langle \sigma\rangle_{inf} &= 0\,,\\
\langle \beta_+\rangle_{inf} &\simeq 0\,,\\
\langle \alpha_+\rangle_{inf} &\simeq -\frac{8g^2\xi\alpha_-}{8g^2\alpha_-{}^2+\lambda^2\alpha_-{}^4-2\lambda^2\alpha_-{}^2+16\lambda^2M^4-8\lambda^2M^2\alpha_-{}^2+8\lambda^2M^2}\,,
\end{aligned} \tag{11}$$

while β_- will represent the main component of the unphysical goldstone boson. Here we use the assumption that $\xi < M_p \ll 2M$. In this case, the effective inflation potential is given by

$$\begin{aligned}
V_{inf}(\alpha_-) &= \frac{1}{16}\left(8g^2\left(\xi-\alpha_-A\right)^2+\lambda^2e^{A^2}\left(\left(A^2-\alpha_-{}^2\right)+4M^2\right)^2\right)\,,\\
A(\alpha_-) &\equiv \frac{8g^2\xi\alpha_-}{\alpha_-{}^2\left(8g^2+\lambda^2\left(\alpha_-{}^2-2\right)\right)+16\lambda^2M^4-8\lambda^2M^2\left(\alpha_-{}^2-1\right)}\,.
\end{aligned} \tag{12}$$

This potential gives the same observables as in the case without FI term in the regime $\xi < M_p \ll M$ [2].

3. CHARGED INFLATON UNDER FLIPPED GUT GAUGE GROUP

In this scenario [3], we study the case when the inflaton is charged under a non-semisimple symmetry groups having the form $\mathcal{G}\times U(1)_X$. It was shown in [14] that such structure is free from producing monopoles when the symmetry is broken at GUT. In this respect, we focus on non-semisimple groups as the unification symmetry, realizing such symmety in two pominent examples: the flipped $SU(5)$ group, $SU(5)\times U(1)$ and the flipped $SO(10)$ group, $SO(10)\times U(1)$ as they preserve the chial structure of the SM. In this case he inflaton will correspond to the righ-handed sneutrino. Again, in order to desribe the model in a suprgravity framework we need to introduce a shift symmetry in a consistent way with the gauge symmetry.

In this case, the Q_X charges are assigned such that the SM hypercharge is given by

$$\begin{aligned}
Y &= \tfrac{1}{5}\left(Q_X-Q_{Y'}\right), & SU(5)\ \ \times U(1)_X,\\
Y &= \tfrac{1}{20}\left(5Q_X-Q_Z-4Q_{Y'}\right), & SO(10)\times U(1)_X,
\end{aligned} \tag{1}$$

where $Q_{Y'}$ is the charge associated with the first abelian factor of the broken $U(1)_{Y'}\times U(1)_X$, subalgebra of $SU(5)\times U(1)_X$ and, Q_Z is the charge associated with the subalgebra $U(1)_{Y'}\times U(1)_Z\times U(1)_X$ of $SO(10)\times U(1)_X$.

Here we will concentrate on the flipped $SU(5)$ group and hence we list the field representations, as well as their charge assignment [3]:

- The standard model (SM) matter content is contained in representations $\mathbf{10}_F$, $\bar{\mathbf{5}}_F$ and $\mathbf{1}_F$, whose respective $U(1)_X$ charges are 1, -3 and 5.
- The Brout-Englert-Higgs bosons enacting electroweak symmetry breaking are contained in $\mathbf{5}_{H_u}$ and $\bar{\mathbf{5}}_{H_d}$.
- The model is complimented with a singlet $\mathbf{1}_S$, necessary to provide the mixing $\mathbf{5}_{H_u}\bar{\mathbf{5}}_{H_d}$ required for electroweak symmetry breaking.
- The heavy scalars Σ and $\bar{\Sigma}$, triggering the breaking of the flipped $SU(5)$ to the standard model gauge group are contained in representations $\mathbf{10}_H$ and $\overline{\mathbf{10}}_H$.
- An additional superfield, in the conjugate representation of the 10-dimensional matter multiplet $\overline{\mathbf{10}}_F$ with $U(1)_X$ charge -1, is added to allow the introduction of a shift symmetry in the Kähler potential, which will be described below.

With these field representations, and taking canonic hypercharge normalization through eq. 1, the generator $Q_{Y'}$ can be written as

$$Q_{Y'} = \frac{1}{6}\mathrm{diag}\left(-2,-2,-2,3,3\right)\,, \tag{2}$$

In addition to this field content, we add Z_2 matter parity in order to forbid undesirable RP-violating couplings [15]. The $U(1)_X$ and Z_2 charges of all the involved fields, can be seen in Table 3.

	$\mathbf{5}_{H_u}$	$\bar{\mathbf{5}}_{H_d}$	$\bar{\mathbf{5}}_F$	$\mathbf{10}_H$	$\overline{\mathbf{10}}_H$	$\mathbf{10}_F$	$\overline{\mathbf{10}}_F$	$\mathbf{1}_F$	$\mathbf{1}_S$
$U(1)_X$	+2	-2	-3	1	-1	1	-1	5	0
Z_2	+	+	-	+	+	-	-	-	+

Here we will focus on the interesting case where the inflaton correspnds to the right-handed sneutrino, N^c ($\bar{N}^c$), embedded in the representation $\mathbf{10}_F$ ($\overline{\mathbf{10}}_F$).

3.1. The inflationary Model

The relevant part of superpotential that is responsible for the inflation scenario is given by [3]

$$W \supset S(\lambda_\phi\, \phi_1\phi_2 + \lambda_h\, h_1 h_2 - M^2) + \mu_\phi \phi_1\phi_2\,, \tag{3}$$

While the shift symmetric Kähler potential is given by [3]

$$K = |\phi_1 + \bar{\phi}_2|^2 + |h_1|^2 + |h_2|^2 + |S|^2 - \eta |S|^4\,, \tag{4}$$

where $\phi_{1(2)} \equiv \overset{(-)}{N^c}$ (contain the inflaton component), $h_{1(2)}$ are the SM singlet component of Σ $(\bar{\Sigma})$, (correspond to the waterfall fields) and S is the stabilizer superfield which is gauge singlet. Accordingly, the supersymmetric vacuum is given by

$$\phi_1 = \phi_2 = S = 0 \;\text{ and }\; h_1 h_2 = \frac{M^2}{\lambda_h}. \tag{5}$$

Therefore the the gauge group $SU(5) \times U(1)$ will be broken at the vacuum and we have $\frac{M^2}{\lambda_h} = M^2_{GUT}$. We perform the following field redefinitions

$$\phi_1 + \bar{\phi}_2 \;=\; \alpha_1 + i\beta_1\,, \qquad \phi_1 - \bar{\phi}_2 = \alpha_2 + i\beta_2\,, \tag{6}$$

$$S \;=\; \frac{s + i\sigma}{\sqrt{2}}\,, \qquad h_{1,2} = \frac{H \pm h}{\sqrt{2}}\,, \tag{7}$$

and finally

$$h \;=\; \frac{h_r + i h_i}{\sqrt{2}}\,, \qquad H = \rho \exp\left(\frac{i}{\sqrt{2}}\frac{\theta}{M_{GUT}}\right)\,. \tag{8}$$

It is clear that θ is the unphysical Goldstone boson and will not contribute to the inflation dynamics. The other spectator fields will be frozen during the inflation as follows

$$\langle\sigma\rangle \;=\; \langle\alpha_1\rangle = \langle\beta_1\rangle = \langle h_r\rangle = \langle h_i\rangle = 0\,, \;\text{ and }\; \langle s\rangle \simeq \frac{\sqrt{2}\,\lambda_\phi\,\mu_\phi}{2\eta\lambda_\phi^2 - \mu_\phi^2}. \tag{9}$$

The waterfall will happen when the field dependent mass squared of the ρ field becomes negative, Assuming that $\mu_\phi \ll \lambda_\phi, M$, and $\lambda_h = \left(\frac{M}{M_{GUT}}\right)^2 \gg M^2$, the mass squared of ρ will be given by

$$m_\rho^2 \;\approx\; \frac{\lambda_\phi^2}{8}\left(\alpha_2^2 + \beta_2^2\right)^2 - \frac{\lambda_\phi\lambda_h}{2}\left(\alpha_2^2 + \beta_2^2\right)\,. \tag{10}$$

While the inflaton rolls down to smaller values, m_ρ^2 becomes negative at the *critical* value

$$\phi_c^2 \equiv \left(\alpha_2^2 + \beta_2^2\right)_c \approx \frac{4\lambda_h}{\lambda_\phi}\,. \tag{11}$$

In [3], it was emphasised that two inflationary regimes can be considered:

- Small critical value: $\phi_c < 1 M_p$

  In this case the FGUT gauge group is broken at the end of inflation and accordingly the effective inflaton potential is given by

  $$V(\alpha_2, \beta_2) = \frac{1}{16}\left(\alpha_2^2 + \beta_2^2\right)\left[(\alpha_2^2 + \beta_2^2)\left(\lambda_\phi^2 - 3\mu_\phi^2\right) + 8(\mu_\phi^2 + \lambda_\phi M^2)\right]. \tag{12}$$

  Due to the rotational symmetry in the plane of α_2 and β_2, the potential can be simplified to the following form

  $$V(\Phi) = \frac{1}{16}\Phi^2\left[\Phi^2\left(\lambda_\phi^2 - 3\mu_\phi^2\right) + 8(\mu_\phi^2 + \lambda_\phi M^2)\right]\,, \tag{13}$$

  with $\Phi = \sqrt{\alpha_2^2 + \beta_2^2}$.

- High critical value: $\phi_c > 20M_p$

  In this case, the field ρ acquires a non-zero vev (an inflaton dependent vev during inflation). This causes a back reaction to the inflation potential but still almost quadratic

$$
\begin{aligned}
V(\Phi) &= \frac{1}{16e\eta^3\lambda_\phi}\Phi^2\Big[8\eta^3\lambda_\phi{\mu_\phi}^2 + 12\eta^3\lambda_h{\mu_\phi}^2 + \eta^2\lambda_\phi\Phi^2\left(\eta\left({\lambda_\phi}^2 - 3{\mu_\phi}^2\right) - {\mu_\phi}^2\right) \\
&+ 12\eta^2\lambda_\phi{\mu_\phi}^2 + 6\eta^2\lambda_\phi{\mu_\phi}^2 + 6\eta\lambda_\phi{\mu_\phi}^2 + \lambda_\phi{\mu_\phi}^2 + 2\eta M^2\left(4\eta^2{\lambda_\phi}^2 + 2\eta{\mu_\phi}^2 + {\mu_\phi}^2\right)\Big].
\end{aligned}
\tag{14}
$$

One can define a relation between μ_ϕ^2 and λ_ϕ^2 in order to see what tuning we need to cancel the effect of the quartic term, as follows

$$
\lambda_\phi^2 = (1+\epsilon)3\mu_\phi^2, \qquad |\epsilon| \ll 1. \tag{15}
$$

Figure 2 shows the inflation observables n_s and r where a scan over the parameters $\mu_\phi = [10^{10}, 10^{12}]$ GeV and $|\epsilon| < 1$ is performed

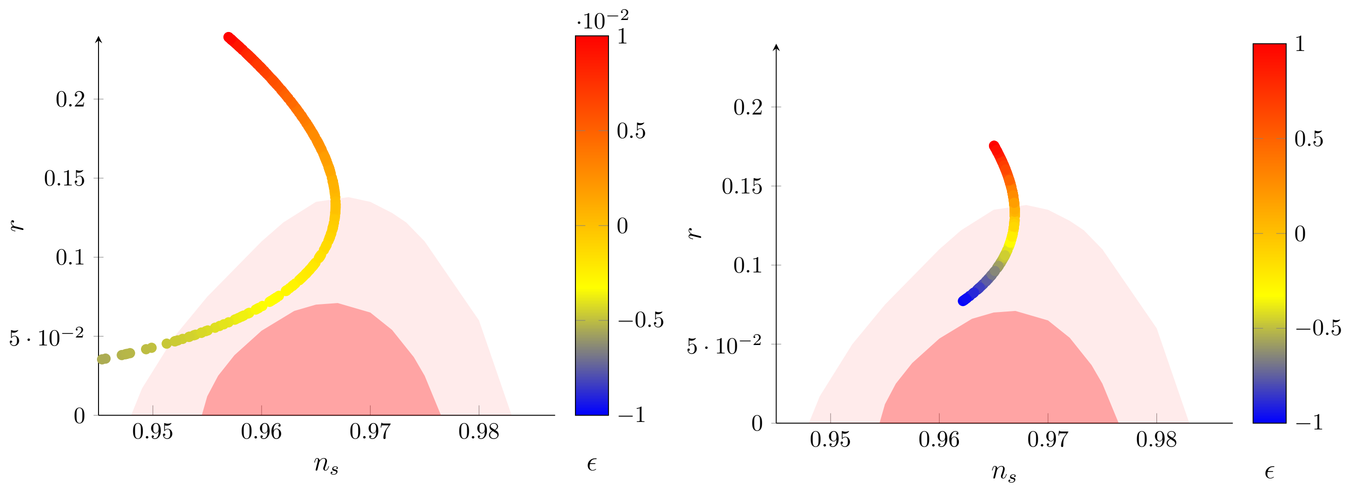

FIGURE 2: Inflation observables for two scenarios where the inflaton field critical values are either small ($\phi_c = 1M_p$, left panel) or large ($\phi_c = 30M_p$, right panel). From Ref. [3].

for the two scenarios of small and large values of ϕ_c. It turns out that the first scenario need a tuning in the value of ϵ to have observables in the 2σ region.

4. REHEATING

The reheating scenario may be one of the interesting feature of such models of inflation. The inflaton rolls down to its minimum then oscillates and decay to the SM particles reheating the universe. In the $U(1)$ charged scenario [2], it was assumed that the inflaton can couple to the right-handed neutrinos by allowing the gauge kinetic mixing between the above inflationary $U(1)$ and the $U(1)_{B-L}$. In this case, the $U(1)_{B-L}$ should be broken at scales $\sim 10^{13}$ GeV (of the same order of the inflaton mass). This is important to have neutrino masses consistent with the observations from one hand and on the other hand to obtain reheating temperature $T_R \sim 10^9$ GeV, consistent with the cosmological constraints [16, 17, 18, 19, 20].

However, in the FGUT scenario [3], it is natural to find couplings between the inflaton and the SM particles from the invariant superpotential under $SU(5) \times U(1)$ as follows

$$
\begin{aligned}
W_{SU(5)} &= Y_u\,\bar{\mathbf{5}}_{H_u\alpha}\,\mathbf{10}_F^{\alpha\beta}\,\bar{\mathbf{5}}_{F\beta} + Y_{d1}\,\epsilon_{\alpha\beta\gamma\delta\lambda}\,\mathbf{10}_F^{\alpha\beta}\,\mathbf{10}_F^{\gamma\delta}\,\mathbf{5}_{H_d}^{\lambda} \\
&+ Y_{d2}\,\epsilon^{\alpha\beta\gamma\delta\lambda}\,\bar{\mathbf{10}}_{F\alpha\beta}\,\bar{\mathbf{10}}_{F\gamma\delta}\,\bar{\mathbf{5}}_{H_u\lambda} + Y_e\,\mathbf{5}_{H_d}^{\lambda}\,\bar{\mathbf{5}}_{F\lambda}\,\mathbf{1}_F,
\end{aligned}
\tag{1}
$$

where the indices $\alpha, \beta, ... = 1, \ldots, 5$. Accordingly, the interaction Lagrangian between the inflaton and the lighter fields of the MSSM can be extracted from the above superpotential as

$$
\mathcal{L}_{int} = -Y\,\tilde{N}^c\left(\nu_L\,\tilde{H}_u^0 + e_L\,\tilde{H}_u^+\right). \tag{2}
$$

The reheating temperature can be calculated using the expression

$$
T_R \approx \frac{(8\pi)^{1/4}}{7}\left(\Gamma\,M_p\right)^{1/2}, \tag{3}
$$

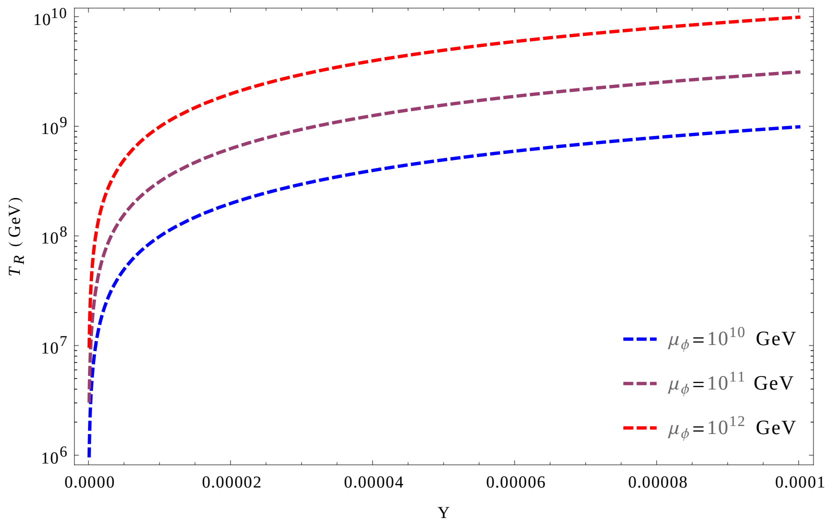

FIGURE 3: Reheating temperature T_R versus the Yukawa coupling Y for different values of the inflaton mass parameter $\mu_\phi = 10^{10}, 10^{11}, 10^{12}$ GeV, from Ref. [3].

with the total decay width of the inflaton Γ is given by

$$\Gamma = \Gamma_{\tilde{\nu}_R \to \nu_L \tilde{H}_u^0} + \Gamma_{\tilde{\nu}_R \to e_L \tilde{H}_u^+}\,, \tag{4}$$

and M_p is the reduced Planck mass. In this respect, the decay of the sneutrino to massless fermions is given by

$$\Gamma_{\tilde{\nu}_R} = \frac{|Y|^2 \mu_\phi}{8\pi}\,. \tag{5}$$

The relation between the reheating temperature and the Yukawa coupling with varying the mass of the inflaton are depicted in Fig. 3. The values of Yukawa couplings the gives consistent reheating temperature are in the range of up-quark coupling value. This indicates that the inflaton comes from the first generation.

5. CONCLUSIONS

In this paper we reviewed a single field inflation scenario within supergravity with shift symmetry where the inflaton is charged under a gauge group. In the case of anomalous and non-anomalous $U(1)$, the effective inflation potential was the same as *new inflation* potential. In the case of FGUT scenario, the effective inflation potential is dominantly quadratic if the coefficient of the quartic term is suppressed. This requires a fine-tuning in the first case where the critical value of the inflaton is small, while the the second case of large critical values is much less tuned. The reheating after inflation due to the decay of the inflaton is analysed in both cases. While it was assumed that the kinetic mixing between our $U(1)$ and $U(1)_{B-L}$ generates couplings between the inflaton and right-handed (s)neutrinos in the first scenario, the second scenario of FGUT gauge symmetry provides natural couplings of the inflaton (the right-handed sneutrino) to the MSSM paricles directly that allows decay channels of the inflaton.

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References

[1] P. A. R. Ade *et al.* [Planck Collaboration], "Planck 2015 results. XX. Constraints on inflation," Astron. Astrophys. **594**, A20 (2016). [arXiv:1502.02114 [astro-ph.CO]].

[2] L. Heurtier, S. Khalil and A. Moursy, "Single Field Inflation in Supergravity with a $U(1)$ Gauge Symmetry," JCAP **1510**, no. 10, 045 (2015). [arXiv:1505.07366 [hep-ph]].

[3] T. E. Gonzalo, L. Heurtier and A. Moursy, "Sneutrino driven GUT Inflation in Supergravity," JHEP **1706**, 109 (2017). [arXiv:1609.09396 [hep-th]].

[4] G. R. Dvali, Q. Shafi and R. K. Schaefer, "Large scale structure and supersymmetric inflation without fine tuning," Phys. Rev. Lett. **73**, 1886 (1994). [hep-ph/9406319].

[5] M. Bastero-Gil, S. F. King and Q. Shafi, "Supersymmetric Hybrid Inflation with Non-Minimal Kahler potential," Phys. Lett. B **651**, 345 (2007) [hep-ph/0604198].

[6] W. Buchmuller, E. Dudas, L. Heurtier and C. Wieck, "Large-Field Inflation and Supersymmetry Breaking," JHEP **1409**, 053 (2014) [arXiv:1407.0253 [hep-th]].
[7] M. Kawasaki, M. Yamaguchi and T. Yanagida, "Natural chaotic inflation in supergravity," Phys. Rev. Lett. **85** (2000) 3572 [hep-ph/0004243].
[8] M. Yamaguchi and J. Yokoyama, "New inflation in supergravity with a chaotic initial condition," Phys. Rev. D **63** (2001) 043506 [hep-ph/0007021].
[9] P. Brax and J. Martin, "Shift symmetry and inflation in supergravity," Phys. Rev. D **72** (2005) 023518 [hep-th/0504168].
[10] R. Kallosh and A. Linde, "New models of chaotic inflation in supergravity," JCAP **1011** (2010) 011 [hep-th/1008.3375].
[11] R. Kallosh, A. Linde and T. Rube, "General inflaton potentials in supergravity," Phys. Rev. D **83** (2011) 043507 [hep-th/1011.5945].
[12] R. Kallosh and A. D. Linde, "Testing String Theory with CMB," JCAP **0704** (2007) 017 [hep-th/0704.0647].
[13] M. Kawasaki and M. Yamaguchi, "A Supersymmetric topological inflation model," Phys. Rev. D **65** (2002) 103518 [hep-ph/0112093].
[14] G. 't Hooft, "Magnetic Monopoles in Unified Gauge Theories," Nucl. Phys. B **79**, 276 (1974).
[15] B. Kyae and Q. Shafi, "Flipped SU(5) predicts delta T/T," Phys. Lett. B **635**, 247 (2006). [hep-ph/0510105].
[16] J. R. Ellis, J. S. Hagelin, D. V. Nanopoulos, K. A. Olive and M. Srednicki, "Supersymmetric Relics from the Big Bang," Nucl. Phys. B **238**, 453 (1984). J. R. Ellis, J. E. Kim and D. V. Nanopoulos, "Cosmological Gravitino Regeneration and Decay," Phys. Lett. B **145**, 181 (1984).
[17] M. Y. Khlopov and A. D. Linde, "Is It Easy to Save the Gravitino?," Phys. Lett. B **138**, 265 (1984).
[18] A. D. Linde, "Particle physics and inflationary cosmology," Contemp. Concepts Phys. **5**, 1 (1990) [hep-th/0503203].
[19] M. Kawasaki and T. Moroi, Prog. Theor. Phys. **93**, 879 (1995) [hep-ph/9403364, hep-ph/9403061].
[20] T. Moroi, [hep-ph/9503210].