**B − L Higgs Inflation in Supergravity With Several Consequences**

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**Abstract:** We consider a renormalizable extension of the minimal supersymmetric standard model endowed by an R and a gauged $B − L$ symmetry. The model incorporates chaotic inflation driven by a quartic potential, associated with the Higgs field which leads to a spontaneous breaking of $U(1)_{B−L}$, and yields possibly detectable gravitational waves. We employ semi-logarithmic Kahler potentials with an enhanced shift symmetry which include only quadratic terms and integer prefactors for the logarithms. An explanation of the $\mu$ term of the MSSM is also provided, consistently with the low energy phenomenology, under the condition that a related parameter in the superpotential is somewhat small. Baryogenesis occurs via non-thermal leptogenesis which is realized by the inflaton’s decay to the lightest and/or next-to-lightest right-handed neutrinos.
1. Introduction

We concentrate on the theoretically most promising models of kinetically modified non-minimal Higgs inflation (HI) investigated in Ref. [1], considering exclusively integer prefactors for the logarithms included in the Kähler potentials. We embed the selected models in a complete framework which presented in Sec. 3. The inflationary part of this context is described in Sec. 3. Then, in Sec. 4, we explain how the minimal supersymmetric standard model (MSSM) is obtained as low energy theory and, in Sec. 5, we outline how the observed baryon asymmetry of the universe (BAU) is generated via non-thermal leptogenesis (nTL). Our conclusions are summarized in Sec. 6. Throughout the text, the subscript of type \( \tilde{z} \) denotes derivation with respect to (w.r.t) the field \( z \) and charge conjugation is denoted by a star. Unless otherwise stated, we use units where \( m_p = 2.433 \cdot 10^{18} \text{ GeV} \) is taken unity.

2. Model Description

We focus on a “Grand Unified Theory” (GUT) based on \( G_{B-L} = G_{SM} \times U(1)_{B-L} \), where \( G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y \) is the gauge group of the standard model and \( B \) and \( L \) denote the baryon and lepton number respectively. The superpotential of our model naturally splits into two parts:

\[
W = W_{MSSM} + W_{HI}, \quad \text{where}
\]

(a) \( W_{MSSM} \) is the part of \( W \) which contains the usual terms – except for the \( \mu \) term – of MSSM, supplemented by Yukawa interactions among the left-handed leptons (\( L_i \)) and \( N_{ij}^c \):

\[
W_{MSSM} = h_{ijD} d_i^c Q_j H_d + h_{ijU} u_i^c Q_j H_u + h_{ijE} e_i^c L_j H_d + h_{ijN} N_{ij}^c L_j H_u. \tag{2.2a}
\]

Here the \( i \)th generation \( SU(2)_L \) doublet left-handed quark and lepton superfields are denoted by \( Q_i \) and \( L_i \) respectively, whereas the \( SU(2)_L \) singlet antiquark [antilepton] superfields by \( u_i^c \) and \( d_i^c \) \( e_i^c \) and \( N_{ij}^c \) respectively. The electroweak Higgs superfields which couple to the up [down] quark superfields are denoted by \( H_u \) \( H_d \).

(b) \( W_{HI} \) is the part of \( W \) which is relevant for HI, the generation of the \( \mu \) term of MSSM and the Majorana masses for \( N_{ij}^c \)’s. It takes the form

\[
W_{HI} = \lambda S (\Phi \Phi - M^2 / 4) + \lambda_{\mu} S H_u H_d + \lambda_{\mu N} \Phi N_{ij}^c \Phi. \tag{2.2b}
\]

The imposed \( U(1)_R \) symmetry ensures the linearity of \( W_{HI} \) w.r.t \( S \). This fact allows us to isolate easily via its derivative the contribution of the inflaton into the F-term SUGRA potential, placing \( S \) at the origin – see Sec. 3.1. It plays also a key role in the resolution of the \( \mu \) problem of MSSM via the second term in the right-hand side (r.h.s) of Eq. (2.2b) – see Sec. 4.2. The inflaton is contained in the system \( \Phi - \Phi \). We are obliged to restrict ourselves to subplanckian values of \( \Phi \) since the imposed symmetries do not forbid non-renormalizable terms of the form \((\Phi \Phi)^p\) with \( p > 1 \) – see Sec. 3.3. The third term in the r.h.s of Eq. (2.2b) provides the Majorana masses for the \( N_{ij}^c \)’s and assures the decay of the inflaton to \( N_{ij}^c \), whose subsequent decay can activate nTL. Here, we work in the so-called \( N_{ij}^c \)-basis, where \( M_{\mu N} \) is diagonal, real and positive. These masses, together with the Dirac neutrino masses in Eq. (2.2a), lead to the light neutrino masses via the seesaw mechanism.
HI is feasible if $W_{HI}$ cooperates with one of the following Kähler potentials:

$$K_1 = -3 \ln \left(1 + c_+ F_+ + F_{1X} |X|^2 \right) + c_- F_- \quad \text{with} \quad F_{1X} = - \ln \left(1 + |X|^2 / N \right),$$  

$$K_2 = -2 \ln \left(1 + c_+ F_+ + c_- F_- + F_{2X} |X|^2 \right) \quad \text{with} \quad F_{2X} = N_X \ln \left(1 + |X|^2 / N_X \right),$$  

$$K_3 = -2 \ln \left(1 + c_+ F_+ + F_{3X} F_- |X|^2 \right) \quad \text{with} \quad F_{3X} = N_X \ln \left(1 + |X|^2 / N_X + c_- F_- / N_X \right),$$

where $F_\pm = |\Phi \pm \bar{\Phi}|^2$, $0 < N_X < 6$, $X^T = S, H_u, H_d, \tilde{N}_i^c$ and the complex scalar components of the superfields $\Phi, \bar{\Phi}, S, H_u, H_d$ and $H_d$ are denoted by the same symbol whereas this of $N_i^c$ by $\tilde{N}_i^c$. The functions $F_\pm$ assist us in the introduction of a shift symmetry for the Higgs fields – cf. Ref. [2]. In all $K$’s, $F_+$ is included in the argument of a logarithm with coefficient $(-3)$ or $(-2)$ whereas $F_-$ is outside it – cf. Ref. [3]. As regards the non-inflaton fields $X^T$, we assume that they have identical kinetic terms expressed by the functions $F_{1X}$ with $l = 1, 2, 3$ and their form is given in Ref. [1]. Just for definiteness, we here adopt the logarithmic forms. These functions ensure the stability and the heaviness and of these modes [4] including exclusively quadratic terms. In the limits $c_+ \to 0$ and $\lambda \to 0$, our models are completely natural in the ’t Hooft sense, since they enjoy the following enhanced symmetries

$$\Phi \to \Phi + c, \quad \bar{\Phi} \to \bar{\Phi} + c^* \quad \text{and} \quad X^T \to e^{i\varphi_\gamma}X^T,$$

where $c$ and $\varphi_\gamma$ are complex and real numbers respectively and no summation is applied over $\gamma$.

### 3. Inflationary Scenario

The salient features of our inflationary scenario are studied at tree level in Sec. [3.1] and at one-loop level in Sec. [3.2]. We then present its predictions in Sec. [3.3], calculating a number of observable quantities introduced in Sec. [3.2].
3.1 Inflationary Potential

If we express $\Phi, \Phi$ and $X^\gamma = S, H_u, H_d, \tilde{N}_i^\gamma$ according to the parametrization

$$\Phi = \phi e^{i\theta} \sqrt{2} \cos \theta_\phi, \quad \Phi = \phi e^{i\theta} \sqrt{2} \sin \theta_\phi \quad \text{and} \quad X^\gamma = x^\gamma + i \bar{x}^\gamma,$$

(3.1)

where $0 \leq \theta_\phi \leq \pi/2$, we can easily deduce that the Einstein frame SUGRA scalar potential $\hat{V}$ which can be found via the formula

$$\hat{V} = \hat{V}_F + \hat{V}_D \quad \text{with} \quad \hat{V}_F = e^K \left(K^{\alpha\beta} D_\alpha W_{H_S} D_\beta W_{H_S} - 3|W_{H_S}|^2\right) \quad \text{and} \quad \hat{V}_D = \frac{1}{2} e^K \sum_a D_a D_a,$$

(3.2)

exhibit a D-flat direction at

$$x^\gamma = \bar{x}^\gamma = \theta = \bar{\theta} = 0 \quad \text{and} \quad \theta_\phi = \pi/4.$$

(3.3)

Along this the only surviving term of $\hat{V}$ can be written universally as

$$\hat{V}_{H_S} = e^K |W_{H_S}|^2 = \frac{\lambda^2 (\phi^2 - M^2)^2}{16 f_R^2} \quad \text{where} \quad f_R = 1 + c_+ \phi^2$$

(3.4)

plays the role of a non-minimal coupling to Ricci scalar in the Jordan frame – see Ref. [2]. Clearly $\hat{V}_{H_S}$ develops an inflationary plateau as in the original case of non-minimal inflation [5]. Contrary to that case, though, here we have also $c_-$ which dominates the canonical normalization of $\phi$ and allows for distinctively different inflationary outputs as shown in Refs. [2, 3]. To specify it together with the normalization of the other fields, we note that, for all $K$’s in Eqs. (2.3a) – (2.3c), $K_{\alpha\beta}$ along the configuration in Eq. (3.3) takes the form

$$\begin{pmatrix} K_{\alpha\beta} \end{pmatrix} = \text{diag} \left( M_{\pm}, K_{\gamma\gamma}, \ldots, K_{\gamma\gamma} \right) \quad \text{with} \quad M_{\pm} = \frac{1}{f_R^2} \begin{pmatrix} \kappa & \bar{\kappa} \\ \bar{\kappa} & \kappa \end{pmatrix} \quad \text{and} \quad K_{\gamma\gamma} = \begin{pmatrix} f_R^{-1} & 1 \\ 1 & f_R \end{pmatrix} \quad \text{for} \quad K = K_1,$$

(3.5)

for $K = K_2, K_3$.

Here $\kappa = c_+ f_R^2 - N c_+$ and $\bar{\kappa} = N c_+^2 \phi^2$ where $N = 3$ [$N = 2$] for $K = K_1$ [$K = K_2$ or $K_3$]. Upon diagonalization of $M_{\pm}$ we find its eigenvalues which are

$$\kappa_+ = c_+ \left(1 + N r_{\pm}(c_+ \phi^2 - 1)/f_R^2\right) \quad \text{and} \quad \kappa_- = c_- \left(1 - N r_{\pm}/f_R\right),$$

(3.6)

where the positivity of $\kappa_-$ is assured during and after HI for

$$r_{\pm} < f_R/N \quad \text{with} \quad r_{\pm} = c_+/c_-.$$  

(3.7)

Given that $f_R > 1$ and $(f_R) \approx 1$, Eq. (3.7) implies that the maximal possible $r_{\pm}$ is $r_{\pm}^{\text{max}} \approx 1/N$. The inequality above discriminates somehow the allowed parameter space for the various choices of $K$’s in Eqs. (2.3a) – (2.3c).

Inserting Eqs. (3.1) and (3.5) into the kinetic term of the SUGRA action, $K_{\alpha\beta} \phi^2 \phi^2$, we can specify the canonically normalized fields, denoted by hat, as follows

$$\frac{d\hat{\phi}}{d\phi} = J, \quad \hat{\theta}_+ = \frac{J}{\sqrt{2}} \phi \theta_+, \quad \hat{\theta}_- = \frac{J}{\sqrt{2}} \phi \theta_-, \quad \hat{\theta}_\phi = \sqrt{\kappa} \phi \left(\theta_\phi - \frac{\pi}{4}\right) \quad \text{and} \quad (\hat{x}^\gamma, \hat{\bar{x}}^\gamma) = \sqrt{\kappa} (x^\gamma, \bar{x}^\gamma),$$

(3.8)

where $J = \sqrt{\kappa_+}$ and $\theta_\pm = (\hat{\theta}_\pm \theta)/\sqrt{2}$. As we show below, the masses of the scalars besides $\hat{\phi}$ during HI are heavy enough such that the dependence of the hatted fields on $\phi$ does not influence their dynamics.
TABLE 2: The mass squared spectrum of our models along the path in Eq. (3.3) for $M \ll \phi \ll 1$ and $N = 3 [N = 2]$ for $K = K_1 [K = K_2$ and $K_3]$. 

3.2 Stability and One-Loop Radiative Corrections

We can verify that the inflationary direction in Eq. (3.3) is stable w.r.t the fluctuations of the non-inflaton fields. To this end, we construct the mass-squared spectrum of the scalars taking into account the canonical normalization of the various fields in Eq. (3.8). In the limit $c_- \gg c_+$, we find the expressions of the masses squared $\tilde{m}^2_{\psi_i}$ (with $\tilde{z}_i = \theta_+, \theta_0, x^3$ and $\tilde{X}^3$) arranged in Table 2. These results approach rather well for $\phi = \phi_*$ – see Sec. 3.3 – the quite lengthy, exact expressions taken into account in our numerical computation. The various unspecified there eigenvalues are defined as follows

$$\tilde{h}_\pm = (\tilde{h}_u \pm \tilde{h}_d) / \sqrt{2}, \quad \tilde{h}_\pm = (\tilde{h}_u \pm \tilde{h}_d) / \sqrt{2} \quad \text{and} \quad \tilde{\psi}_\pm = (\tilde{\psi}_{\phi_\pm} \pm \tilde{\psi}_3) / \sqrt{2}, \quad (3.9a)$$

where the (unhatted) spinors $\psi_\phi$ and $\psi_\Phi$ associated with the superfields $\Phi$ and $\bar{\Phi}$ are related to the normalized (hatted) ones in Table 2 as follows

$$\tilde{\psi}_{\phi_\pm} = \sqrt{K_{\phi}} \psi_{\phi_\pm} \quad \text{with} \quad \psi_{\phi_\pm} = (\psi_{\phi_\pm} \pm \psi_3) / \sqrt{2}. \quad (3.9b)$$

From Table 2 it is evident that $0 < N_X \leq 6$ assists us to achieve $m^2_{\tilde{h}_\pm} > \tilde{H}^2_{\tilde{h}} = \tilde{V}_{\tilde{h}} / 3$ – in accordance with the results of Ref. [4] – and also enhances the ratios $m^2_{\tilde{X}^3}/\tilde{H}^2_{\tilde{h}}$ for $X^3 = H_u, H_d, N_3$ w.r.t the values that we would have obtained, if we had used just canonical terms in the $K$’s. On the other hand, $m^2_{\tilde{h}_\pm} > 0$ requires

$$\lambda_\mu < \lambda (1 + c_+ \phi^2 / 2) / 4 (1 / \phi^2 + c_+) \quad \text{for} \quad K = K_1; \quad (3.10a)$$

$$\lambda_\mu < \lambda \phi^2 (1 + 1 / N_X) / 4 \quad \text{for} \quad K = K_2 \text{ and } K_3. \quad (3.10b)$$

In both cases, the quantity in the r.h.s of the inequality takes its minimal value at $\phi = \phi_t$ – see Sec. 3.3 – and numerically equals to $2 \cdot 10^{-5} - 10^{-6}$. In Table 2 we display also the mass $M_{BL}$ of the gauge boson which becomes massive having ‘eaten’ the Goldstone boson $\theta_\pm$. This signals the fact that $G_{B-L}$ is broken during HI and so no cosmological defects are produced. Also, we can verify [1] that radiative corrections à la Coleman-Weinberg can be kept under control.
3.3 Inflationary Observables

A period of slow-roll HI is controlled by the strength of the slow-roll parameters

\[ \hat{\varepsilon} = \frac{1}{2} \left( \frac{\dot{V}_{\text{HI}}}{V_{\text{HI}}} \right)^2 \approx \frac{8}{c_- \phi^2 \mathcal{f}_R} \quad \text{and} \quad \hat{\eta} = \frac{\ddot{V}_{\text{HI}}}{V_{\text{HI}}} \approx \frac{12}{c_- \phi^2 \mathcal{f}_R^2}. \]  

(3.11)

Expanding \( \hat{\varepsilon} \) and \( \hat{\eta} \) for \( \phi \ll 1 \) we can find that HI terminates for \( \phi = \phi_t \) such that

\[ \max \{ \hat{\varepsilon}(\phi_t), |\hat{\eta}(\phi_t)| \} = 1 \quad \Rightarrow \quad \phi_t \simeq \max \left\{ \frac{2\sqrt{2/c_-}}{\sqrt{1+16r_\pm}}, \frac{2\sqrt{3/c_-}}{\sqrt{1+36r_\pm}} \right\}. \]  

(3.12)

The number of e-foldings, \( \hat{N}_* \), that the pivot scale \( k_* = 0.05/\text{Mpc} \) suffers during HI can be calculated through the relation

\[ \hat{N}_* = \int_{\phi_t}^{\phi_*} d\phi \frac{\dot{V}_{\text{HI}}}{V_{\text{HI}}} \simeq \frac{1}{16r_\pm} \left( (1 + c_+ \phi^2_*)^2 - 1 \right). \]  

(3.13)

where \( \phi_* [\phi_*] \) is the value of \( \hat{\phi} [\phi] \) when \( k_* \) crosses the inflationary horizon. Given that \( \phi_t \ll \phi_* \), we can write \( \phi_* \) as a function of \( \hat{N}_* \) as follows

\[ \phi_* \simeq \sqrt{(f_{R_*} - 1)/c_+} \quad \text{with} \quad f_{R_*} = \left( 1 + 16r_\pm \hat{N}_* \right)^{1/2}. \]  

(3.14)

We can impose a lower bound on \( c_- \) above which \( \phi_* \leq 1 \) for every \( r_\pm \). Indeed, from Eq. (3.14) we have

\[ \phi_* \leq 1 \quad \Rightarrow \quad c_- \geq (f_{R_*} - 1)/r_\pm. \]  

(3.15)

and so, our proposal can be stabilized against corrections from higher order terms of the form \((\Phi \Phi)^p\) with \( p \geq 1 \) in \( W_{\text{HI}} \) – see Eq. (2.22). Despite the fact that \( c_- \) may take relatively large values, the corresponding effective theory is valid up to \( m_p = 1 \). To clarify further this point we have to identify the ultraviolet cut-off scale \( \Lambda_{\text{UV}} \) of theory analyzing the small-field behavior of our models. More specifically, we expand about \( \langle \phi \rangle = M \ll 1 \) the kinetic term \( J^2 \dot{\phi}^2 \) in the SUGRA action [1] and \( \dot{V}_{\text{HI}} \) in Eq. (3.3). Our results can be written in terms of \( \hat{\phi} \) as

\[ J^2 \dot{\phi}^2 \simeq \left( 1 + 3N r_\pm^2 \hat{\phi}^2 - 5N r_\pm^2 \hat{\phi}^4 + \cdots \right) \hat{\phi}^2 \quad \text{and} \quad \dot{V}_{\text{HI}} \simeq \frac{\lambda}{16c_-} \left( 1 - 2r_\pm \hat{\phi}^2 + 3r_\pm^2 \hat{\phi}^4 - \cdots \right). \]  

(3.16)

From the expressions above we conclude that \( \Lambda_{\text{UV}} = m_p \) since \( r_\pm \leq 1 \) due to Eq. (3.7).

The power spectrum \( A_s \) of the curvature perturbations generated by \( \phi \) at the pivot scale \( k_* \) is estimated as follows

\[ \sqrt{A_s} = \frac{1}{2\sqrt{3\pi}} \frac{\dot{V}_{\text{HI}}(\hat{\phi}_*)^{3/2}}{|\dot{V}_{\text{HI}}(\hat{\phi}_*)|} \simeq \frac{\lambda}{32\sqrt{3\pi}} \phi^3_\ast \quad \Rightarrow \quad \lambda = 32\sqrt{3A_s \pi c_-} \left( \frac{r_\pm}{f_{R_*} - 1} \right)^{3/2}. \]  

(3.17)

The resulting relation reveals that \( \lambda \) is proportional to \( c_- \) for fixed \( r_\pm \).

At the same pivot scale, we can also calculate \( n_s \), its running, \( a_s \), and \( r \) via the relations

\[ n_s = 1 - 6\hat{\varepsilon}_* + 2\hat{\eta}_* \simeq 1 - \frac{3}{2N_*} - \frac{3}{8(N_*^2 r_\pm)^{1/2}}, \quad r = 16\hat{\varepsilon}_* \simeq \frac{1}{2N_*} \frac{2}{(N_*^2 r_\pm)^{1/2}}. \]  

(3.18a)

\[ a_s = \frac{2}{3} \left( 4\hat{\eta}_*^2 - (n_s - 1)^2 \right) - 2\hat{\xi}_* \simeq - \frac{3}{2N_*^2} \quad \text{with} \quad \hat{\xi} = \frac{\dot{V}_{\text{HI}}}{\dot{V}_{\text{HI}}(\hat{\phi})^{3/2}}/V_{\text{HI}}. \]  

(3.18b)

Here the variables with subscript * are evaluated at \( \phi = \phi_* \). A clear dependence of \( n_s \) and \( r \) on \( r_\pm \) arises.
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9.55 9.60 9.65 9.70 9.75 9.80
1 2 3 4 5 6 7 8 9 10
0.025 0.015 0.06 0.5
9 \times 10^{-3}

9.10^{-3} 0.015 0.025 0.06 0.5

0.002 (0.01)
ns (0.1)

-3

4.8 \times 10^{-3}

\lambda = 6.3 \times 10^{-3}, r_e = 0.025

\phi_\star (a)

\phi_f (b)

V_{HI}^{(10^{-9})}

\phi_\star \approx 61.5 + \ln \frac{\hat{V}_{HI}(\phi_\star)^{1/2}}{\hat{V}_{HI}(\phi_\star)^{1/4}} + \frac{1}{2} f_R(\phi_\star)

\lambda = 6.3 \times 10^{-3}, r_e = 0.025

\hat{N}_s \approx 4.627 \times 10^{-5},

\hat{V}_{HI}^{(10^{-3})}

\phi_\star

F I G U R E 1: (a) Allowed curve in the $n_s - r_{0.002}$ plane for $K = K_2$ or $K_3$ with the $r_{\pm}$ values indicated on it – the marginalized joint 68% [95%] regions from Planck, BAO and BK14 data [7] are depicted by the dark [light] shaded contours; (b) The inflationary potential $\hat{V}_{HI}$ as a function of $\phi$ for $\phi > 0$, $r_{\pm} \simeq 0.025$ and $\lambda = 6.3 \times 10^{-3}$ – the values of $\phi_\star, \phi_f$ and are also indicated.

3.4 COMPARISON WITH OBSERVATIONS

The approximate analytic expressions above can be verified by the numerical analysis of our model. Namely, we apply the accurate expressions in Eqs. (3.13) and (3.17) and confront the corresponding observables with the requirements [5, 6]

\begin{align}
(a) \hat{N}_s & \approx 61.5 + \ln \frac{\hat{V}_{HI}(\phi_\star)^{1/2}}{\hat{V}_{HI}(\phi_\star)^{1/4}} + \frac{1}{2} f_R(\phi_\star) & \text{and} \quad (b) A_s^{1/2} & \approx 4.627 \times 10^{-5},
\end{align}

where we consider in Eq. (3.13a) an equation-of-state parameter $w_{int} = 1/3$ corresponding to quartic potential which is expected to approximate rather well $\hat{V}_{HI}$ for $\phi \ll 1$. We, thus, restrict $\lambda$ and $\phi_\star$ and compute the model predictions via Eqs. (3.13a) and (3.13b) for any selected $r_{\pm}$. These must be in agreement with the fitting of the Planck, Baryon Acoustic Oscillations (BAO) and BICEP2/Keck Array data [5, 7] with $\Lambda$CDM+$r$ model, i.e.,

\begin{align}
(a) n_s & = 0.968 \pm 0.009 \quad \text{and} \quad (b) r \leq 0.07,
\end{align}

at 95% confidence level (c.l.) with $|a_s| \ll 0.01$.

Let us clarify here that the free parameters of our models are $r_{\pm}$ and $\lambda/c_-$ and not $c_+, c_-$ and $\lambda$ as naively expected. Indeed, if we perform the rescalings

\begin{align}
\Phi & \to \Phi/\sqrt{c_-}, \quad \bar{\Phi} \to \bar{\Phi}/\sqrt{c_-} \quad \text{and} \quad S \to S,
\end{align}

$W$ in Eq. (2.2b) depends on $\lambda/c_-$ and the $K$’s in Eq. (2.3a) – (2.3c) depend on $r_{\pm}$. As a consequence, $\hat{V}_{HI}$ depends exclusively on $\lambda/c_-$ and $r_{\pm}$. Since the $\lambda/c_-$ variation is rather trivial – see Ref. [3] – we focus on the variation of the other parameters.

Our results are displayed in Fig. 1 for $K = K_2$ or $K_3$. Namely, in Fig. 1-(a) we show a comparison of the models’ predictions against the observational data [5, 7] in the $n_s - r_{0.002}$ plane, where $r_{0.002} =$
16\hat{\hat{c}}(\phi_{0.002}) with \hat{\phi}_{0.002} being the value of \hat{\phi} when the scale k = 0.002/Mpc, which undergoes \hat{N}_{0.002} = \hat{N}_x + 3.22 e-foldings during HI, crosses the horizon of HI. We depict the theoretically allowed values with a solid line with the variation of r_± shown along it. For low enough r_±’s – i.e. r_± ≤ 0.0005 – the line reaches (n_s, r_{0.002}) ≃ (0.947, 0.28) obtained within minimal quartic inflation defined for c_+ = 0. Increasing r_± the line enters the observationally allowed regions and terminates for r_± ≃ 0.5, beyond which Eq. (3.7) is violated. Along this line we find – consistently with the analytic formulas of Sec. 3.3

\[ 0.048 \lesssim r_\pm \lesssim 0.1, \quad 9.64 \lesssim \frac{n_s}{0.1} \lesssim 9.72, \quad 0.7 \lesssim \frac{r}{0.01} \lesssim 8.1 \quad \text{and} \quad 0.17 \lesssim \frac{10^5 \lambda}{c_-} \lesssim 3.13. \]  

Moreover a_s ≃ -(5 - 6) \cdot 10^{-4} and so, our models are consistent with the fitting of data with the ΛCDM+r model [5]. These are also testable by the forthcoming experiments, like BICEP3, PRISM and LiteBIRD, searching for primordial gravity waves since r ≥ 0.007. Had we employed K = K1, the line in Fig. 4.(a) would have been shortened until r_± ≃ 0.33 yielding r_{0.002} ≥ 0.0084. The other bounds would have been remained more or less unaffected.

Taking the χ^2 distribution of the obtained (n_s, r)'s we can identify the following best-fit value:

\[ r_\pm = 0.025 \quad \text{resulting to} \quad (n_s, r) = (0.969, 0.033). \]  

For this value we display the structure of \hat{V}_{HI} as a function of \hat{\phi} in Fig. 1(b). We take φ_s = 1 which corresponds to \lambda = 6.3 \cdot 10^{-3} and c_- = 146. We observe that \hat{V}_{HI} is a monotonically increasing function of \phi. The inflationary scale, \hat{V}_{HI}^{1/4}, approaches the SUSY GUT scale \hat{M}_{GUT} ≃ 8.2 \cdot 10^{-3} and lies well below \hat{A}_{UV} = 1, consistently with the classical approximation to the inflationary dynamics.

4. HIGGS INFLATION AND μ TERM OF MSSM

A byproduct of the R symmetry associated with our model is that it assists us to understand the origin of μ term of MSSM, as we show in Sec. 4.1, consistently with the low-energy phenomenology of MSSM – see Sec. 4.3. Hereafter we restore units, i.e., we take m_p = 2.433 \cdot 10^{18} \text{ GeV}.

4.1 SUSY POTENTIAL

The SUSY limit \hat{V}_{SUSY} of \hat{V}_{HI} in Eq. (3.4) is given by

\[ V_{SUSY} = \hat{K}^\alpha_\beta W^{\alpha}_{HI} W^{\beta}_{HI} + \frac{g^2}{2} \sum D_\alpha D_\alpha, \]  

where \hat{K} is the limit of the K’s in Eqs. (2.3a) – (2.3d) for m_p → ∞. Focusing on the S – \Phi – \Phi system we find

\[ \hat{K} = c_- F_- - N c_+ F_+ + |S|^2. \]  

Upon substitution of \hat{K} into Eq. (4.1a) we obtain

\[ V_{SUSY} = \lambda^2 \left| \Phi \right|^2 \left| \Phi - \frac{1}{4} M^2 \right|^2 + \frac{\lambda^2}{c_- (1 - N r_\pm)} S^2 \left( \left| \Phi \right|^2 + \left| \Phi \right|^2 \right) + \frac{g^2}{2} c_-^2 (1 - N r_\pm)^2 \left( \left| \Phi \right|^2 - \left| \Phi \right|^2 \right)^2. \]  

From the last equation, we find that the SUSY vacuum lies along the D-flat direction |\Phi| = |\Phi| with

\[ \langle S \rangle = 0 \quad \text{and} \quad |\langle \Phi \rangle| = |\langle \Phi \rangle| = M/2. \]  

As a consequence, ⟨\Phi⟩ and ⟨\bar{\Phi}⟩ break spontaneously U(1)_{B-L} down to \mathbb{Z}_2^{B-L}. Since U(1)_{B-L} is already broken during HI, no cosmic string are formed.
4.2 Generation of the $\mu$ Term of MSSM

The contributions from the soft SUSY breaking terms, although negligible during HI, since these are much smaller than $\phi$, may shift slightly $\langle S \rangle$ from zero in Eq. (4.2). Indeed, the relevant potential terms are

$$V_{\text{soft}} = \left( \lambda A_{\lambda} \bar{S} \Phi + \lambda \mu A_{\lambda} \bar{S} H d + \lambda N' A_{\lambda} \bar{N}' - a_3 S \bar{\lambda} M^2/4 + \text{h.c.} \right) + m_2^2 |X|^2,$$

(4.3)

where $m_r, A_{\lambda}, A_{\mu}, A_{N'}$, and $a_3$ are soft SUSY breaking mass parameters. Rotating $S$ in the real axis by an appropriate $R$-transformation, choosing conveniently the phases of $A_{\lambda}$ and $a_3$ so as the total low energy potential $V_{\text{tot}} = V_{\text{SUSY}} + V_{\text{soft}}$ to be minimized – see Eq. (4.1) – and substituting in $V_{\text{soft}}$ the SUSY vacuum expectation values (v.e.v.s) of $\Phi$ and $\bar{\Phi}$ from Eq. (4.2) we get

$$\langle V_{\text{tot}}(S) \rangle = \lambda^2 M^2 S^2/2 c_-(1 - N r_\pm) - \lambda a_3/2 m_{3/2}^2 M^2 S,$$

(4.4a)

where we take into account that $m_S \ll M$ and we set $|A_{\lambda}| = |a_3| = 2a_3/2 m_{3/2}$ with $m_{3/2}$ being the $\tilde{G}$ mass and $a_{3/2} > 0$ a parameter of order unity which parameterizes our ignorance for the dependence of $|A_{\lambda}|$ and $|a_3|$ on $m_{3/2}$. The minimization condition for the total potential in Eq. (4.4a) w.r.t $S$ leads to a non vanishing $\langle S \rangle$ as follows

$$\frac{d}{dS} \langle V_{\text{tot}}(S) \rangle = 0 \Rightarrow \langle S \rangle \simeq a_3/2 m_{3/2} c_-/(1 - N r_\pm)/\lambda.$$

(4.4b)

The generated $\mu$ term from the second term in the r.h.s of Eq. (2.2c) is

$$\mu = \lambda_\mu \langle S \rangle \simeq \lambda_\mu a_3/2 m_{3/2} c_- (1 - N r_\pm)/\lambda.$$

(4.5)

By virtue of Eq. (3.17), the resulting $\mu$ above depends on $r_\pm$ and does not depend on $\lambda$ and $c_-$. We may verify that any $|\mu|$ value is accessible for the $\lambda_\mu$ values allowed by Eqs. (3.10a) and (3.10b) without any ugly hierarchy between $m_{3/2}$ and $\mu$.

4.3 Connection with the MSSM Phenomenology

The SUSY breaking effects, considered in Eq. (4.3), explicitly break $U(1)_R$ to a subgroup, $\mathbb{Z}_2^K$ which can be identified with a matter parity. Under this discrete symmetry all the matter (quark and lepton) superfields change sign – see Table 1. From the $R$ charges there we conclude that $\mathbb{Z}_2^K$ remains unbroken, since $\langle S \rangle$ in Eq. (4.4b) also breaks spontaneously $U(1)_R$ to $\mathbb{Z}_2^K$ and so no disastrous domain walls are formed. Combining $\mathbb{Z}_2^K$ with the $\mathbb{Z}_2^F$ fermion parity, under which all fermions change sign, yields the well-known $R$-parity. This residual symmetry prevents rapid proton decay and guarantees the stability of the lightest SUSY particle (LSP), providing thereby a well-motivated cold dark matter (CDM) candidate.

The candidacy of LSP may be successful, if it generates the correct CDM abundance [6] within a concrete low energy framework. In our case this is the MSSM or, more specifically, the Constrained MSSM (CMSSM), if we adopt only the following free parameters

$$\text{sign} \mu, \tan \beta = \langle H_u \rangle/\langle H_d \rangle, M_{1/2}, m_0, \text{ and } A_0,$$

(4.6)

where $\text{sign} \mu$ is the sign of $\mu$, and the three last mass parameters denote the common gaugino mass, scalar mass and trilinear coupling constant, respectively, defined (normally) at $M_{\text{GUT}}$. The parameter
| CMSSM REGION | $|A_0|$ (TeV) | $m_0$ (TeV) | $|\mu|$ (TeV) | $a_{3/2}$ | $\lambda_\mu$ ($10^{-6}$) |
|-------------|-------------|-------------|-------------|--------|------------------|
| $A/H$ Funnel | 9.9244      | 1.36        | 1.409       | 1.086  | 0.6223           |
| $\tilde{\tau}_1 - \chi$ Coannihilation | 1.2271      | 1.476       | 2.62        | 0.831  | 9.36             |
| $\tilde{\tau}_1 - \chi$ Coannihilation | 9.965       | 4.269       | 4.073       | 2.33   | 1.794            |
| $\tilde{\chi}_1^+ - \chi$ Coannihilation | 9.2061      | 9.000       | 0.983       | 1.023  | 0.468            |

$\mu$ is not free, since it is computed at low scale by enforcing the conditions for the electroweak symmetry breaking. The values of the (four and one half) parameters in Eq. (4.6) can be tightly restricted imposing a number of cosmo-phenomenological constraints from which the consistency of LSP relic density with observations plays a central role. Some updated results are recently presented in Ref. [8], where we can also find the best-fit values of $|A_0|$, $m_0$ and $|\mu|$ listed in the first four lines of Table 3.

We see that there are four allowed regions characterized by the specific mechanism for suppressing the relic density of the LSP which is the lightest neutralino ($\chi$) – $\tilde{\tau}_1$, $\tilde{\tau}_1$, and $\tilde{\chi}_1^+$ stand for the lightest stau, stop and chargino eigenstate. If we take the best-fit value of $r_{\pm}$ in Eq. (3.23) and identify

$$m_0 = m_{3/2} \text{ and } |A_0| = |A_\chi| = |a_5|$$

we can derive first $a_{3/2}$ and then the $\lambda_\mu$ values which yield the phenomenologically desired $|\mu|$ – ignoring renormalization group effects. The outputs of our computation is listed in the two rightmost columns of Table 3 for $K = K_1$, $K_2$ and $K_3$. From these we infer that the required $\lambda_\mu$ values, in all cases besides the one, written in italics, are comfortably compatible with Eqs. (3.10a) and (3.10b) for $N_X = 2$ which imply

$$\lambda_\mu \lesssim 6.6 \cdot 10^{-6} \text{ for } K = K_1 \text{ and } \lambda_\mu \lesssim 1.1 \cdot 10^{-5} \text{ for } K = K_2 \text{ and } K_3.$$  \hspace{1cm} (4.8)

Therefore, we conclude that the whole inflationary scenario can be successfully combined with all the allowed regions CMSSM besides the $\tilde{\tau}_1 - \chi$ coannihilation region for $K = K_1$. On the other hand, all the CMSSM regions can be consistent with the gravitino limit on $T_{\text{th}}$ – see Sec. 5.2. Indeed, $m_{3/2}$ as low as 1 TeV becomes cosmologically safe, under the assumption of the unstable $\tilde{G}$, for the $T_{\text{th}}$ values, necessitated for satisfactory leptogenesis, as presented in Table 4.

5. Non-Thermal Leptogenesis and Neutrino Masses

We below specify how our inflationary scenario makes a transition to the radiation dominated era (Sec. 5.1) and offers an explanation of the observed BAU (Sec. 5.2) consistently with the $\tilde{G}$ constraint and the low energy neutrino data. Our results are summarized in Sec. 5.3.
5.1 Inflaton Mass & Decay

When HI is over, the inflaton continues to roll down towards the SUSY vacuum, Eq. (4.2). Soon after, it settles into a phase of damped oscillations around the minimum of $\tilde{V}_{	ext{HI}}$. The (canonically normalized) inflaton,

$$\tilde{\delta}\phi = \langle J \rangle \delta\phi \quad \text{with} \quad \delta\phi = \phi - M \quad \text{and} \quad \langle J \rangle = \sqrt{\langle \kappa_+ \rangle} = \sqrt{c_- (1 - N r_\pm)} \quad (5.1)$$

acquires mass, at the SUSY vacuum in Eq. (4.3), which is given by

$$\tilde{m}_{\delta\phi} = \sqrt{\tilde{V}_{\text{HI}} / \langle J \rangle^2} = \left( \frac{\tilde{V}_{\text{HI},\tilde{\phi}} / \langle J \rangle}{\delta\phi} \right)^{1/2} \approx \frac{\lambda M}{\sqrt{c_- (1 - N r_\pm)}} \quad (5.2)$$

where the last (approximate) equality above is valid only for $r_\pm \ll 1/N$ – see Eqs. (3.8) and (5.8). As we see, $\tilde{m}_{\delta\phi}$ depends crucially on $M$ which may be, in principle, a free parameter acquiring any subplanckian value without disturbing the inflationary process. To determine better our models, though, we prefer to specify $M$ requiring that $\langle \Phi \rangle$ and $\langle \tilde{\Phi} \rangle$ in Eq. (4.2) take the values dictated by the unification of the MSSM gauge coupling constants, despite the fact that $U(1)_{\text{B-L}}$ gauge symmetry does not disturb this unification and $M$ could be much lower. In particular, the unification scale $M_{\text{GUT}} \simeq 2 \cdot 10^{16}$ GeV can be identified with $M_{\text{BL}}$ – see Table 2 – at the SUSY vacuum in Eq. (4.2), i.e.,

$$\frac{\sqrt{c_- (f_R) - N r_\pm}}{g \langle f_R \rangle} \equiv M_{\text{GUT}} \Rightarrow M \simeq M_{\text{GUT}} / g \sqrt{c_- (1 - N r_\pm)} \quad (5.3)$$

with $g \simeq 0.7$ being the value of the GUT gauge coupling and we take into account that $\langle f_R \rangle \simeq 1$. Upon substitution of the last expression in Eq. (5.3) into Eq. (5.2) we can infer that $\tilde{m}_{\delta\phi}$ remains constant for fixed $r_\pm$ since $\lambda / c_-$ is fixed too – see Eq. (5.18). Particularly, along the line in Fig. 1-(a) we obtain

$$3.5 \cdot 10^{11} \lesssim \tilde{m}_{\delta\phi} / \text{GeV} \lesssim 3.9 \cdot 10^{13} \quad \text{for} \quad K = K_1; \quad (5.4a)$$

$$3.46 \cdot 10^{10} \lesssim \tilde{m}_{\delta\phi} / \text{GeV} \lesssim 4.2 \cdot 10^{13} \quad \text{for} \quad K = K_2 \quad \text{and} \quad K_3; \quad (5.4b)$$

During the phase of its oscillations at the SUSY vacuum, $\tilde{\delta}\phi$ decays perturbatively reheating the Universe at a reheat temperature given by

$$T_{\text{rh}} = \left( \frac{72}{5 \pi^2 g_*} \right)^{1/4} \left( \frac{\phi}{m_{\text{pl}}} \right)^{1/2} \quad \text{with} \quad \phi = \tilde{\phi}_{\text{HI}} \rightarrow N_i \tilde{N}_i \rightarrow H_u H_d. \quad (5.5)$$

Also $g_* = 228.75$ counts the MSSM effective number of relativistic degrees of freedom. To compute $T_{\text{rh}}$ we take into account the following decay widths:

$$\Gamma_{\delta\phi \rightarrow N_i \tilde{N}_i} = \frac{\tilde{m}_{\delta\phi}^2}{16 \pi} \left( 1 - \frac{4 M_{\text{IN}}^2}{\tilde{m}_{\delta\phi}^2} \right)^{3/2} \quad \text{with} \quad g_{\text{IN}} = \frac{\lambda_{\text{IN}}}{\langle J \rangle} \left( 1 - 3 c_+ N \frac{M^2}{m_{\text{pl}}^2} \right) \quad (5.6a)$$

$$\Gamma_{\delta\phi \rightarrow H_u H_d} = \frac{\tilde{m}_{\delta\phi}}{8 \pi} \frac{g_{H}}{m_{\text{pl}}} \quad \text{with} \quad g_{H} = \frac{\lambda_{\text{H}}}{\sqrt{2}} \left( 1 - 2 c_+ \frac{M^2}{m_{\text{pl}}^2} \right) \quad (5.6b)$$

arising from the lagrangian terms

$$\mathcal{L}_{\delta\phi \rightarrow N_i \tilde{N}_i} = - \frac{1}{2} K / \phi^2 W_{HI,\tilde{N}_i \tilde{N}_i^c N_i} + \text{h.c.} = g_{\text{IN}} \tilde{\delta}\phi (N_i^c \tilde{N}_i + \text{h.c.}) + \cdots, \quad (5.6c)$$

$$\mathcal{L}_{\delta\phi \rightarrow H_u H_d} = - \tilde{\delta}\phi \hat{K}^{SS} |W_{HI,S}|^2 = - g_{H} \tilde{m}_{\delta\phi} \tilde{\delta}\phi (H_u^c H_d + \text{h.c.}) + \cdots \quad (5.6d)$$
describing $\tilde{\phi}$ decay into a pair of $N_i^c$ with Majorana masses $M_{jN^c} = \hat{\lambda}_{iN^c} M$ and $H_u$ and $H_d$ respectively. Note that the decay modes into MSSM (s)-particles $XYZ$ [1] through a typical trilinear superpotential term of MSSM is suppressed since they arise from non-renormalizable interactions proportional to $M/m_T \ll 1$.

### 5.2 Lepton-Number and Gravitino Abundances

For $T_{\text{th}} < M_{N^c}$, the out-of-equilibrium decay of $N_i^c$ generates a lepton-number asymmetry (per $N_i^c$ decay), $\varepsilon_i$. The resulting lepton-number asymmetry is partially converted through sphaleron effects into a yield of the observed BAU:

$$Y_B = -0.35 \cdot \frac{5}{2} \frac{T_{\text{th}}}{m_{\tilde{\phi}}} \sum \frac{\Gamma_{\tilde{\phi} \rightarrow N_i^c N_j^c}}{\Gamma_{\tilde{\phi}}} \varepsilon_i,$$

with $\varepsilon_i = \sum_{j} \frac{\text{Im} \left[ (m_{D}^i m_{D}^j)^2 \right]}{8 \pi (H_u)^2 (m_{D}^i m_{D}^j)^2} \left( F_{\delta} (x_{ij}, y_i, y_j) + F_{v} (x_{ij}) \right)$, (5.7)

where $\langle H_u \rangle \simeq 174$ GeV, for large $\tan \beta$, $m_0$ the Dirac mass matrix of $\nu_i$ and $F_{\delta}$ [$F_{v}$] are the functions entered in the vertex and self-energy contributions computed as indicated in Ref. [9]. The expression above has to reproduce the observational result [6]:

$$Y_B = (8.64^{+0.15}_{-0.16}) \cdot 10^{-11}.$$  

(5.8)

The validity of Eq. (5.7) requires that the $\tilde{\phi}$ decay into a pair of $N_i^c$'s is kinematically allowed for at least one species of the $N_i^c$'s and also that there is no erasure of the produced $Y_L$ due to $N_i^c$ mediated inverse decays and $\Delta L = 1$ scatterings. These prerequisites are ensured if we impose

$$\text{(a) } \hat{m}_{\tilde{\phi}} \geq 2M_{N^c} \text{ and (b) } M_{1N^c} \gtrsim 10 T_{\text{th}},$$

(5.9)

The quantity $\varepsilon_i$ can be expressed in terms of the Dirac masses of $\nu_i$, $m_{D}$, arising from the third term of Eq. (5.12). Employing the seesaw formula we can then obtain the light-neutrino mass matrix $m_{\nu}$ in terms of $m_{D}$ and $M_{N^c}$. As a consequence, nTL can be nicely linked to low energy neutrino data. We take as inputs the recently updated best-fit values [10] – cf. Ref. [1] – on the neutrino mass-squared differences, $\Delta m^2_{12} = 7.65 \cdot 10^{-5}$ eV$^2$ and $\Delta m^2_{21} = 2.55 \cdot 10^{-3}$ eV$^2$ [\Delta m^2_{12} = 2.49 \cdot 10^{-3}$ eV$^2$], on the mixing angles, $\sin^2 \theta_{12} = 0.321$, $\sin^2 \theta_{13} = 0.02155$ [\sin^2 \theta_{13} = 0.0214] and $\sin^2 \theta_{23} = 0.43$ [\sin^2 \theta_{23} = 0.596] and the CP-violating Dirac phase $\delta = 1.4\pi \ [\delta = 1.4\pi] \text{ for normal [inverted] ordered (NO [IO]) neutrino masses, } m_{\nu}$'s. Furthermore, the sum of $m_{\nu}$'s is bounded from above at 95% c.l. by the data [6],

$$\sum m_{\nu} \leq 0.23 \text{ eV.}$$

(5.10)

The required $T_{\text{th}}$ in Eq. (5.7) must be compatible with constraints on the gravitino ($\tilde{G}$) abundance, $Y_{3/2}$, at the onset of nucleosynthesis (BBN), which is estimated to be:

$$Y_{3/2} \simeq 1.9 \cdot 10^{-22} \frac{T_{\text{th}}}{\text{GeV}},$$

(5.11)

where we take into account only thermal production of $\tilde{G}$, and assume that $\tilde{G}$ is much heavier than the MSSM gauginos. On the other hand, $Y_{3/2}$ is bounded from above in order to avoid spoiling the success of the BBN. For the typical case where $G$ decays with a tiny hadronic branching ratio, we have

$$Y_{3/2} \lesssim \begin{cases} 10^{-14} & \text{for } m_{3/2} \simeq 0.69 \text{ TeV,} \\
10^{-13} & \text{implying } T_{\text{th}} \lesssim 5.3 \cdot 10^7 \text{ GeV,} \\
10^{-12} & \text{ } 10^8 \text{ GeV,} \\
10^{-11} & \text{ } 10^9 \text{ GeV.} \end{cases}$$

(5.12)

The bounds above can be somehow relaxed in the case of a stable $\tilde{G}$. 

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**B-L HI in SUGRA With Several Consequences**

**CONSTANTINOS PALLIS**
| PARAMETERS | CASES |
|-----------|-------|
|           | A     | B     | C     | D     | E     | F     | G     |
|           | Normal | Almost | Hierarchy | Degeneracy | Inverted | Hierarchy |
| Low Scale Parameters |
| $m_{1\nu}/0.1\text{ eV}$ | 0.05 | 0.1 | 0.5 | 0.7 | 0.7 | 0.5 | 0.49 |
| $m_{2\nu}/0.1\text{ eV}$ | 0.1 | 0.13 | 0.51 | 0.7 | 0.7 | 0.51 | 0.5 |
| $m_{3\nu}/0.1\text{ eV}$ | 0.5 | 0.51 | 0.7 | 0.86 | 0.5 | 0.1 | 0.05 |
| $\sum m_{\nu}/0.1\text{ eV}$ | 0.65 | 0.74 | 1.7 | 2.3 | 1.9 | 1.1 | 1 |
| $\varphi_1$ | $-\pi/5$ | $-\pi/2$ | $\pi$ | $\pi/9$ | 0 | $-3\pi/4$ | $\pi/2$ |
| $\varphi_2$ | $\pi$ | 0 | $\pi/3$ | $\pi$ | $\pi/2$ | 5$\pi/4$ | $-\pi/2$ |
| Leptogenesis-Scale Parameters |
| $m_{11D}/0.1\text{ GeV}$ | 2 | 5 | 10 | 10 | 1.3 | 5 | 6 |
| $m_{21D}/\text{GeV}$ | 6 | 1.97 | 3.9 | 10 | 9 | 0.715 | 1.1 |
| $m_{31D}/\text{GeV}$ | 100 | 150 | 170 | 168 | 202 | 100 | 199 |
| $M_{1N^c}/10^{10}\text{ GeV}$ | 1.0 | 3.3 | 2.85 | 3.3 | 2.98 | 0.45 | 1.23 |
| $M_{2N^c}/10^{10}\text{ GeV}$ | 6.9 | 13.6 | 26.5 | 111.4 | 13.9 | 2.76 | 2.8 |
| $M_{3N^c}/10^{14}\text{ GeV}$ | 2.9 | 4.9 | 2.2 | 1.2 | 3.7 | 2.7 | 27.2 |
| Open Decay Channels of the Inflaton, $\tilde{\delta}\phi$, into $N^c_i$ |
| $\tilde{\delta}\phi \rightarrow$ | $N^c_1$ | $N^c_1$ | $N^c_1$ | $N^c_1$ | $N^c_{1,2}$ | $N^c_{1,2}$ |
| $\tilde{\Gamma}_{\tilde{\delta}\phi \rightarrow N^c_i N^c_j}/\tilde{\Gamma}_{\tilde{\delta}\phi}$ (%) | 13.8 | 15.4 | 17.4 | 14.9 | 17.1 | 18.3 | 22.7 |
| Resulting B-Yield |
| $10^{11}Y_B$ | 8.68 | 8.66 | 8.79 | 8.69 | 8.58 | 8.67 | 8.68 |
| Resulting $T_{th}$ and $G$-Yield |
| $T_{th}/10^7\text{ GeV}$ | 2.8 | 2.8 | 2.84 | 2.8 | 2.84 | 2.85 | 2.94 |
| $10^{15}Y_{3/2}$ | 5.3 | 5.3 | 5.4 | 5.3 | 5.4 | 5.4 | 5.5 |

**Table 4**: Parameters yielding the correct $Y_B$ for various neutrino mass schemes. We take $K = K_2$ or $K_3$ with $N_x = 2$, $r_{\mu} = 0.025$ and $\lambda_{\mu} = 10^{-6}$.

**5.3 Results**

Confronting with observations $Y_B$ and $Y_{3/2}$ which depend on $\hat{m}_{\delta\phi}$, $T_{th}$, $M_{IN^c}$ and $m_{1D}$’s – see Eqs. (5.7) and (5.11) – we can further constrain the parameter space of the our models. We follow
the bottom-up approach detailed in Ref. [1], according to which we find the $M_{\nu^N}$’s by using as inputs the $m_{123}$’s, a reference mass of the $V_{ij}$’s – $m_{11}$’s for NO $m_{123}$’s, or $m_{33}$’s for IO $m_{123}$’s –, the two Majorana phases $\phi_1$ and $\phi_2$ of the PMNS matrix, and the best-fit values for the low energy parameters of neutrino physics mentioned in Sec. 3.2. In our numerical code, we also estimate [1] the RG evolved values of the latter parameters at the scale of nTL, $\Lambda_{10}$, by considering the MSSM with $\tan \beta \simeq 50$ as an effective theory between $\Lambda_{10}$ and the soft SUSY breaking scale, $M_{\text{SUSY}} = 1.5$ TeV. We evaluate the $M_{\nu^N}$’s at $\Lambda_{10}$, and we neglect any possible running of the $m_{123}$’s and $M_{\nu^N}$’s. The so obtained $M_{\nu^N}$’s clearly correspond to the scale $\Lambda_{10}$.

Some representative values of the parameters which yield $Y_B$ and $Y_{3/2}$ compatible with Eqs. (5.8) and (5.12), respectively are arranged in Table 4. We take the best-fit $r_\pm$ value in Eq. (5.23) and $\lambda_\mu = 10^{-6}$ in accordance with Eqs. (6.10a) and (3.10a) with $N_X = 2$. We obtain $\hat{m}_{\delta \beta} = 8.6 \cdot 10^{10} \text{ GeV}$ for $K = K_1$ and $\hat{m}_{\delta \phi} = 8.6 \cdot 10^{10} \text{ GeV}$ for $K = K_2$ or $K_3$. Although such an uncertainty from the choice of $K$’s do not cause any essential alteration of the final outputs, we mention just for definiteness that we take $K = K_2$ or $K_3$ throughout. We consider NO (cases A and B), almost degenerate (cases C, D and E) and IO (cases F and G) $m_{123}$’s. In all cases, the current limit in Eq. (5.10) is safely met. The gauge symmetry considered here does not predict any particular Yukawa unification pattern and so, the $m_{123}$’s are free parameters. This fact facilitates the fulfilment of Eq. (5.9b) since $mD[1]$ affects heavily $mrh[1]$. Care is also taken so that the perturbativity of $\lambda_{\nu^N}$ holds, i.e., $\lambda_{\nu^N}^2/4\pi \leq 1$. The inflaton $\widetilde{\delta \phi}$ decays mostly into $N_i^\nu$’s – see cases A – E. In all cases $\Gamma_{\delta \phi \rightarrow N_i^\nu N_j^\nu} < \Gamma_{\delta \phi \rightarrow H_u H_d}$ and so the ratios $\Gamma_{\delta \phi \rightarrow N_i^\nu N_j^\nu}/\Gamma_{\delta \phi \rightarrow H_u H_d}$ introduce a considerable reduction in the derivation of $Y_B$. In Table 4 we also display the values of $T_{h_h}$, the majority of which are close to $3 \cdot 10^7 \text{ GeV}$, and the corresponding $Y_{3/2}$’s, which are consistent with Eq. (5.12) for $m_{3/2} \geq 1 \text{ TeV}$. These values are in nice agreement with the ones needed for the solution of the $\mu$ problem of MSSM – see, e.g., Table 5.

In order to investigate the robustness of the conclusions inferred from Table 4, we examine also how the central value of $Y_B$ in Eq. (5.8) can be achieved by varying $r_\pm$ (or $\hat{m}_{\delta \beta}$) and adjusting conveniently $m_{2D}$ (or $M_{2\nu}$) – see Fig. 2(a) or (b) respectively. We fix again $\lambda_\mu = 10^{-6}$. Since the range
of $Y_B$ in Eq. (5.8) is very narrow, the 95% c.l. width of these contours is negligible. The convention adopted for the various lines is depicted in the plot of Fig. 3(b). In particular, we use solid, dashed and dot-dashed line when the remaining inputs – i.e. $m_{1\nu}$, $m_{1D}$, $m_{3D}$, $\phi_1$, and $\phi_2$ – correspond to the cases B, C and F of Table 3, respectively. Only some segments from the $r_\pm$ range in Eq. (3.22) fulfill the post-inflationary requirements. Namely, as inferred by Fig. 2-(a), we find that $r_\pm$ may vary in the ranges $(0.0161 - 0.18)$, $(0.025 - 0.21)$ and $(0.025 - 0.499)$ for $m_{2D}$ plotted in Fig. 2-(a) and the remaining inputs of the cases B, C and F respectively. As regards the other quantities, in all we obtain

$$4.4 \lesssim Y_G/10^{-15} \lesssim 228 \quad \text{and} \quad 0.23 \lesssim T_{th}/10^8\text{GeV} \lesssim 12 \quad \text{with} \quad 6.5 \lesssim \tilde{m}_{\tilde{g}_8}/10^{10} \lesssim 4241. \quad \text{(5.13)}$$

As a bottom line, nTL not only is a realistic possibility within our setting but also it can be comfortably reconciled with the $\tilde{G}$ constraint even for $m_3/2 \sim 1\text{ TeV}$ as deduced from Eqs. (5.13) and (5.12).

6. CONCLUSIONS

We investigated the realization of kinetically modified non-minimal HI and nTL in the framework of a $B - L$ extension of MSSM endowed with the condition that the GUT scale is determined by the running of the three gauge coupling constants. Our setup is tied to the super- and Kähler potentials given in Eqs. (2.2b) and (2.3a) – (2.3c). Prominent in this setting is the role of a softly broken shift-symmetry whose violation is parameterized by the quantity $r_\pm = c_+/c_-$ and can be constrained by the observations. Our models exhibit the following features: (i) they inflate away cosmological defects; (ii) they safely accommodates observable gravitational waves with subplanckian inflaton values and without causing any problem with the validity of the effective theory; (iii) they offer a nice solution to the $\mu$ problem of MSSM, provided that $\lambda_\mu$ is somehow small; (iv) they allow for baryogenesis via nTL compatible with $\tilde{G}$ constraints and neutrino data. In particular, we may have $m_{3/2} \sim 1\text{ TeV}$, with the inflaton decaying mainly to $N_c1$ and $N_c2$ – we obtain $M_{\tilde{N}_c}$ in the range $(10^9 - 10^{14})$ GeV. It remains to introduce a consistent soft SUSY breaking sector in the theory which is certainly an important and difficult task.

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