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$B-L$ Higgs Inflation in Supergravity with Several Consequences

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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 μ 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.

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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 presented in Sec. 2. 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 $,z$ 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 $m_{\rm P} = 2.433\cdot 10^{18}$ GeV is taken unity.

2. Model Description

We focus on a "*Grand Unified Theory*" (GUT) based on $G_{B-L} = G_{\rm SM}\times U(1)_{B-L}$, where $G_{\rm SM} = SU(3)_{\rm C}\times SU(2)_{\rm L}\times U(1)_Y$ 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:

$$W = W_{\rm MSSM} + W_{\rm HI},\ \text{ where} \tag{2.1}$$

(a) $W_{\rm MSSM}$ is the part of W which contains the usual terms – except for the μ term – of MSSM, supplemented by Yukawa interactions among the left-handed leptons (L_i) and N^c_i:

$$W_{\rm MSSM} = h_{ijD}d^c_i Q_j H_d + h_{ijU}u^c_i Q_j H_u + h_{ijE}e^c_i L_j H_d + h_{ijN}N^c_i L_j H_u. \tag{2.2a}$$

Here the ith generation $SU(2)_{\rm L}$ doublet left-handed quark and lepton superfields are denoted by Q_i and L_i respectively, whereas the $SU(2)_{\rm L}$ singlet antiquark [antilepton] superfields by u^c_i and ${d_i}^c$ [e^c_i and N^c_i] respectively. The electroweak Higgs superfields which couple to the up [down] quark superfields are denoted by H_u [H_d].

(b) $W_{\rm HI}$ is the part of W which is relevant for HI, the generation of the μ term of MSSM and the Majorana masses for N^c_i's. It takes the form

$$W_{\rm HI} = \lambda S\left(\bar\Phi\Phi - M^2/4\right) + \lambda_\mu S H_u H_d + \lambda_{iN^c}\bar\Phi N^{c2}_i\,. \tag{2.2b}$$

The imposed $U(1)_R$ symmetry ensures the linearity of $W_{\rm HI}$ 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 μ 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 $\bar\Phi-\Phi$. We are obliged to restrict ourselves to subplanckian values of $\bar\Phi\Phi$ since the imposed symmetries do not forbid non-renormalizable terms of the form $(\bar\Phi\Phi)^p$ 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^c_i's and assures the decay of the inflaton to $\widetilde N^c_i$, whose subsequent decay can activate nTL. Here, we work in the so-called N^c_i*-basis*, where M_{iN^c} 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.

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SUPERFIELDS	REPRESENTATIONS UNDER G_{B-L}	GLOBAL SYMMETRIES		
		R	B	L
MATTER FIELDS				
e_i^c	$(\mathbf{1},\mathbf{1},1,1)$	1	0	-1
N_i^c	$(\mathbf{1},\mathbf{1},0,1)$	1	0	-1
L_i	$(\mathbf{1},\mathbf{2},-1/2,-1)$	1	0	1
u_i^c	$(\mathbf{3},\mathbf{1},-2/3,-1/3)$	1	$-1/3$	0
d_i^c	$(\mathbf{3},\mathbf{1},1/3,-1/3)$	1	$-1/3$	0
Q_i	$(\bar{\mathbf{3}},\mathbf{2},1/6,1/3)$	1	$1/3$	0
HIGGS FIELDS				
H_d	$(\mathbf{1},\mathbf{2},-1/2,0)$	0	0	0
H_u	$(\mathbf{1},\mathbf{2},1/2,0)$	0	0	0
S	$(\mathbf{1},\mathbf{1},0,0)$	2	0	0
Φ	$(\mathbf{1},\mathbf{1},0,2)$	0	0	-2
$\bar{\Phi}$	$(\mathbf{1},\mathbf{1},0,-2)$	0	0	2

Table 1: *The representations under G_{B-L} and the extra global charges of the superfields of our model.*

HI is feasible if $W_{\rm HI}$ cooperates with *one* of the following Kähler potentials:

$$K_1 = -3\ln\left(1+c_+F_+ + F_{1X}(|X|^2)\right) + c_-F_- \quad \text{with} \quad F_{1X} = -\ln\left(1+|X|^2/N\right), \tag{2.3a}$$

$$K_2 = -2\ln\left(1+c_+F_+\right) + c_-F_- + F_{2X}(|X|^2) \quad \text{with} \quad F_{2X} = N_X\ln\left(1+|X|^2/N_X\right), \tag{2.3b}$$

$$K_3 = -2\ln\left(1+c_+F_+\right) + F_{3X}(F_-,|X|^2) \quad \text{with} \quad F_{3X} = N_X\ln\left(1+|X|^2/N_X + c_-F_-/N_X\right), \tag{2.3c}$$

where $F_\pm = \left|\Phi\pm\bar{\Phi}^*\right|^2$, $0<N_X<6$, $X^\gamma = S, H_u, H_d, \widetilde{N}_i^c$ and the complex scalar components of the superfields $\Phi, \bar{\Phi}, S, H_u$ and H_d are denoted by the same symbol whereas this of N_i^c by $\widetilde{N}_i^c$. 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^γ, we assume that they have identical kinetic terms expressed by the functions F_{lX} with $l=1,2,3$ and their form is given in Ref. [1]. Just for definiteness, we here adopt the logarithmic forms. These functions ensures 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

$$\Phi\to\ \Phi+c,\ \ \bar{\Phi}\to\ \bar{\Phi}+c^* \ \text{ and }\ X^\gamma\to\ e^{i\varphi_\gamma}X^\gamma, \tag{2.4}$$

where c and φ_γ are complex and real numbers respectively and no summation is applied over γ.

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. 3.4, calculating a number of observable quantities introduced in Sec. 3.3.

3.1 Inflationary Potential

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

$$\Phi = \frac{\phi e^{i\theta}}{\sqrt{2}}\cos\theta_\Phi, \ \ \bar{\Phi} = \frac{\phi e^{i\bar{\theta}}}{\sqrt{2}}\sin\theta_\Phi \ \text{ and } \ X^\gamma = \frac{x^\gamma + i\bar{x}^\gamma}{\sqrt{2}}, \tag{3.1}$$

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

$$\widehat{V} = \widehat{V}_{\rm F} + \widehat{V}_{\rm D} \ \text{ with } \ \widehat{V}_{\rm F} = e^K\left(K^{\alpha\bar{\beta}} D_\alpha W_{\rm HI} D^*_{\bar{\beta}} W^*_{\rm HI} - 3|W_{\rm HI}|^2\right) \ \text{ and } \ \widehat{V}_{\rm D} = \frac{1}{2} g^2 \sum_{\rm a} {\rm D_a D_a}, \tag{3.2}$$

exhibit a D-flat direction at

$$x^\gamma = \bar{x}^\gamma = \theta = \bar{\theta} = 0 \ \text{ and } \ \theta_\Phi = \pi/4. \tag{3.3}$$

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

$$\widehat{V}_{\rm HI} = e^K K^{SS^*} |W_{{\rm HI},S}|^2 = \frac{\lambda^2(\phi^2 - M^2)^2}{16 f_R^2} \ \text{ where } \ f_R = 1 + c_+\phi^2 \tag{3.4}$$

plays the role of a non-minimal coupling to Ricci scalar in the *Jordan frame* – see Ref. [2]. Clearly $\widehat{V}_{\rm HI}$ 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 ϕ 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), $K_{\alpha\bar{\beta}}$ along the configuration in Eq. (3.3) takes the form

$$\left(K_{\alpha\bar{\beta}}\right) = \mathsf{diag}\left(M_\pm, \underbrace{K_{\gamma\bar{\gamma}}, ..., K_{\gamma\bar{\gamma}}}_{8 \text{ elements}}\right) \ \text{ with } \ M_\pm = \frac{1}{f_R^2}\begin{pmatrix} \kappa & \bar{\kappa} \\ \bar{\kappa} & \kappa \end{pmatrix} \ \text{ and } \ K_{\gamma\bar{\gamma}} = \begin{cases} f_R^{-1} & \text{for } K = K_1, \\ 1 & \text{for } K = K_2, \ K_3. \end{cases} \tag{3.5}$$

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

$$\kappa_+ = c_-\left(1 + N r_\pm (c_+\phi^2 - 1)/f_R^2\right) \simeq c_- \ \text{ and } \ \kappa_- = c_-\left(1 - N r_\pm / f_R\right), \tag{3.6}$$

where the positivity of κ_- is assured during and after HI for

$$r_\pm < f_R/N \ \text{ with } \ r_\pm = c_+/c_-. \tag{3.7}$$

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

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

$$\frac{d\widehat{\phi}}{d\phi} = J, \ \ \widehat{\theta}_+ = \frac{J}{\sqrt{2}}\phi\theta_+, \ \ \widehat{\theta}_- = \sqrt{\frac{\kappa_-}{2}}\phi\theta_-, \ \ \widehat{\theta}_\Phi = \sqrt{\kappa_-}\phi\left(\theta_\Phi - \frac{\pi}{4}\right) \ \text{ and } \ (\widehat{x}^\gamma, \widehat{\bar{x}}^\gamma) = \sqrt{K_{\gamma\bar{\gamma}}}(x^\gamma, \bar{x}^\gamma), \tag{3.8}$$

where $J = \sqrt{\kappa_+}$ and $\theta_\pm = \left(\bar{\theta} \pm \theta\right)/\sqrt{2}$. As we show below, the masses of the scalars besides $\widehat{\phi}$ during HI are heavy enough such that the dependence of the hatted fields on ϕ does not influence their dynamics.

EIGEN-STATES	MASSES SQUARED			
		$K=K_1$	$K=K_2$	$K=K_3$
$\widehat{\theta}_+$	$\widehat{m}^2_{\theta+}$	$6\widehat{H}^2_{\rm HI}$		$6(1+1/N_X)\widehat{H}^2_{\rm HI}$
$\widehat{\theta}_\Phi$	$\widehat{m}^2_{\theta_\Phi}$	$M^2_{BL}+\widehat{m}^2_{\theta+}$		
$\widehat{s},\widehat{\bar{s}}$	$\widehat{m}^2_s$	$6c_+\phi^2\widehat{H}^2_{\rm HI}/N$	$6\widehat{H}^2_{\rm HI}/N_X$	
$\widehat{h}_\pm,\widehat{\bar{h}}_\pm$	$\widehat{m}^2_{h\pm}$	$3\widehat{H}^2_{\rm HI}\left(1\pm4\lambda_\mu(\phi^{-2}+c_+)/\lambda\right)+\widehat{m}^2_s/2$	$3\widehat{H}^2_{\rm HI}\left(1+1/N_X\pm4\lambda_\mu/\lambda\phi^2\right)$	
$\widehat{\tilde{\nu}}^c_i,\widehat{\bar{\tilde{\nu}}}^c_i$	$\widehat{m}^2_{i\tilde{\nu}^c}$	$3\widehat{H}^2_{\rm HI}\left(1+16\lambda^2_{iN^c}(\phi^{-2}+c_+)/\lambda^2\right)+\widehat{m}^2_s/2$	$3\widehat{H}^2_{\rm HI}\left(1+1/N_X+16\lambda^2_{iN^c}/\lambda^2\phi^2\right)$	
A_{BL}	M^2_{BL}	$g^2c_-\left(1-Nr_\pm/f_R\right)\phi^2$		
$\widehat{\psi}_\pm$	$\widehat{m}^2_{\psi\pm}$	$6\left((N-3)c_+\phi^2-2\right)^2\widehat{H}^2_{\rm HI}/c_-\phi^2f^2_R$	$6\left((N-2)c_+\phi^2-2\right)^2\widehat{H}^2_{\rm HI}/c_-\phi^2f^2_R$	
$\widehat{N}^c_i$	$\widehat{m}^2_{iN^c}$	$48\lambda^2_{iN^c}\widehat{H}^2_{\rm HI}/\lambda^2\phi^2$		
$\lambda_{BL},\widehat{\psi}_{\Phi-}$	M^2_{BL}	$g^2c_-\left(1-Nr_\pm/f_R\right)\phi^2$		

Table 2: *The mass squared spectrum of our models along the path in Eq. (3.3) for $M\ll\phi\ll1$ 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 $\widehat{m}^2_{z^\alpha}$ (with $z^\alpha=\theta_+,\theta_\Phi,x^\gamma$ and $\bar{x}^\gamma$) arranged in Table 2. These results approach rather well for $\phi=\phi_\star$ – 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

$$\widehat{h}_\pm=(\widehat{h}_u\pm\widehat{h}_d)/\sqrt{2},\quad \widehat{\bar{h}}_\pm=(\widehat{\bar{h}}_u\pm\widehat{\bar{h}}_d)/\sqrt{2}\quad\text{and}\quad \widehat{\psi}_\pm=(\widehat{\psi}_{\Phi+}\pm\widehat{\psi}_S)/\sqrt{2}, \tag{3.9a}$$

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

$$\widehat{\psi}_{\Phi\pm}=\sqrt{\kappa_\pm}\psi_{\Phi\pm}\quad\text{with}\quad \psi_{\Phi\pm}=(\psi_\Phi\pm\psi_{\bar{\Phi}})/\sqrt{2}. \tag{3.9b}$$

From Table 2 it is evident that $0<N_X\leq6$ assists us to achieve $m^2_s>\widehat{H}^2_{\rm HI}=\widehat{V}_{\rm HI}/3$ – in accordance with the results of Ref. [4] – and also enhances the ratios $m^2_{X^{\tilde{\gamma}}}/\widehat{H}^2_{\rm HI}$ for $X^{\tilde{\gamma}}=H_u,H_d,\widetilde{N}^c_i$ 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, $\widehat{m}^2_{h-}>0$ requires

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

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

In both cases, the quantity in the r.h.s of the inequality takes its minimal value at $\phi=\phi_{\rm f}$ – see Sec. 3.3 – and numerically equals to $2\cdot10^{-5}-5\cdot10^{-6}$. In Table 2 we display also the mass M_{BL} of the gauge boson which becomes massive having 'eaten' the Goldstone boson θ_-. 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 á 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

$$\widehat{\varepsilon}=\frac{1}{2}\left(\frac{\widehat{V}_{\rm HI,\widehat{\phi}}}{\widehat{V}_{\rm HI}}\right)^2\simeq\frac{8}{c_-\phi^2 f_R^2}\quad\text{and}\quad\widehat{\eta}=\frac{\widehat{V}_{\rm HI,\widehat{\phi}\widehat{\phi}}}{\widehat{V}_{\rm HI}}\simeq 12\,\frac{1-c_+\phi^2}{c_-\phi^2 f_R^2}\,\cdot \tag{3.11}$$

Expanding $\widehat{\varepsilon}$ and $\widehat{\eta}$ for $\phi\ll 1$ we can find that HI terminates for $\phi=\phi_{\rm f}$ such that

$$\max\{\widehat{\varepsilon}(\phi_{\rm f}),|\widehat{\eta}(\phi_{\rm f})|\}=1\quad\Rightarrow\quad\phi_{\rm f}\simeq\max\left\{\frac{2\sqrt{2/c_-}}{\sqrt{1+16r_\pm}},\frac{2\sqrt{3/c_-}}{\sqrt{1+36r_\pm}}\right\}. \tag{3.12}$$

The number of e-foldings, $\widehat{N}_\star$, that the pivot scale $k_\star=0.05/{\rm Mpc}$ suffers during HI can be calculated through the relation

$$\widehat{N}_\star=\int_{\widehat{\phi}_{\rm f}}^{\widehat{\phi}_\star}d\widehat{\phi}\,\frac{\widehat{V}_{\rm HI}}{\widehat{V}_{\rm HI,\widehat{\phi}}}\simeq\frac{1}{16r_\pm}\left((1+c_+\phi_\star^2)^2-1\right), \tag{3.13}$$

where $\widehat{\phi}_\star$ $[\phi_\star]$ is the value of $\widehat{\phi}$ $[\phi]$ when $k_\star$ crosses the inflationary horizon. Given that $\phi_{\rm f}\ll\phi_\star$, we can write $\phi_\star$ as a function of $\widehat{N}_\star$ as follows

$$\phi_\star\simeq\sqrt{(f_{R\star}-1)/c_+}\ \text{ with }\ f_{R\star}=\left(1+16r_\pm\widehat{N}_\star\right)^{1/2}. \tag{3.14}$$

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

$$\phi_\star\leq 1\quad\Rightarrow\quad c_-\geq(f_{R\star}-1)/r_\pm \tag{3.15}$$

and so, our proposal can be stabilized against corrections from higher order terms of the form $(\Phi\bar\Phi)^p$ with $p>1$ in $W_{\rm HI}$ – see Eq. (2.2b). Despite the fact that c_- may take relatively large values, the corresponding effective theory is valid up to $m_{\rm P}=1$. To clarify further this point we have to identify the ultraviolet cut-off scale $\Lambda_{\rm UV}$ 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 $\widehat{V}_{\rm HI}$ in Eq. (3.4). Our results can be written in terms of $\widehat{\phi}$ as

$$J^2\dot\phi^2\simeq\left(1+3Nr_\pm^2\widehat{\phi}^2-5Nr_\pm^3\widehat{\phi}^4+\cdots\right)\dot{\widehat{\phi}}^2\quad\text{and}\quad\widehat{V}_{\rm HI}\simeq\frac{\lambda^2\widehat{\phi}^4}{16c_-^2}\left(1-2r_\pm\widehat{\phi}^2+3r_\pm^2\widehat{\phi}^4-\cdots\right). \tag{3.16}$$

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

The power spectrum $A_{\rm s}$ of the curvature perturbations generated by ϕ at the pivot scale $k_\star$ is estimated as follows

$$\sqrt{A_{\rm s}}=\frac{1}{2\sqrt{3}\,\pi}\,\frac{\widehat{V}_{\rm HI}(\widehat{\phi}_\star)^{3/2}}{|\widehat{V}_{\rm HI,\widehat{\phi}}(\widehat{\phi}_\star)|}\simeq\frac{\lambda\sqrt{c_-}}{32\sqrt{3}\pi}\phi_\star^3\quad\Rightarrow\quad\lambda=32\sqrt{3A_{\rm s}}\pi c_-\left(\frac{r_\pm}{f_{R\star}-1}\right)^{3/2}\cdot \tag{3.17}$$

The resulting relation reveals that λ is proportional to c_- for fixed $r_\pm$.

At the same pivot scale, we can also calculate $n_{\rm s}$, its running, $a_{\rm s}$, and r via the relations

$$n_{\rm s}=1-6\widehat{\varepsilon}_\star+2\widehat{\eta}_\star\simeq 1-\frac{3}{2\widehat{N}_\star}-\frac{3}{8(\widehat{N}_\star^3 r_\pm)^{1/2}},\quad r=16\widehat{\varepsilon}_\star\simeq\frac{1}{2\widehat{N}_\star^2 r_\pm}+\frac{2}{(\widehat{N}_\star^3 r_\pm)^{1/2}}, \tag{3.18a}$$

$$a_{\rm s}=\frac{2}{3}\left(4\widehat{\eta}_\star^2-(n_{\rm s}-1)^2\right)-2\widehat{\xi}_\star\simeq-\frac{3}{2\widehat{N}_\star^2}\ \text{ with }\ \widehat{\xi}=\widehat{V}_{\rm HI,\widehat{\phi}}\widehat{V}_{\rm HI,\widehat{\phi}\widehat{\phi}\widehat{\phi}}/\widehat{V}_{\rm HI}^2\,. \tag{3.18b}$$

Here the variables with subscript $\star$ are evaluated at $\phi=\phi_\star$. A clear dependence of $n_{\rm s}$ and r on $r_\pm$ arises.

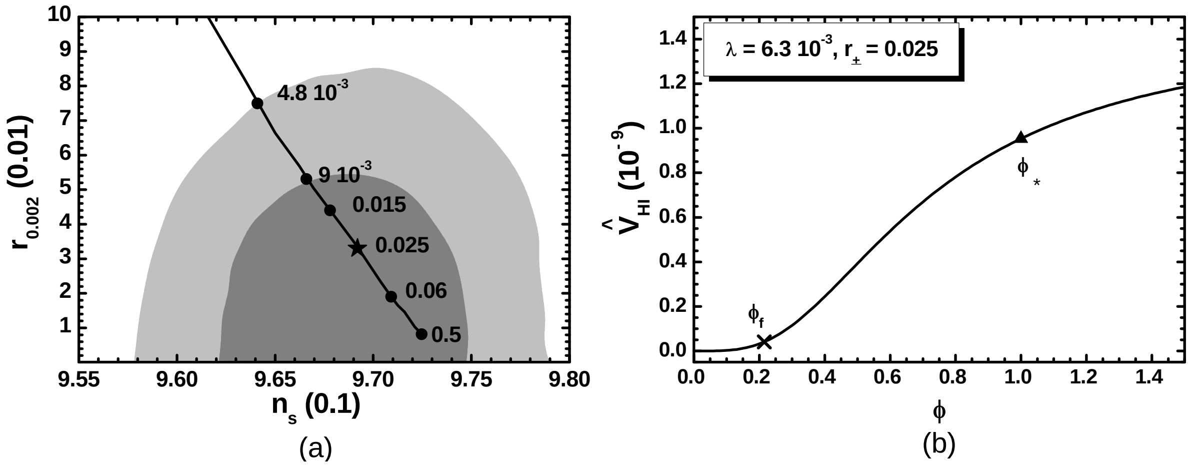

Figure 1: *(a) Allowed curve in the $n_s - r_{0.002}$ plane for $K = K_2$ or K_3 with the $r_\pm$ values indicated on it – the marginalized joint 68% [95%] regions from Planck, BAO and BK14 data [7] are depicted by the dark [light] shaded contours; (b) The inflationary potential $\widehat{V}_{\rm HI}$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

$$\text{(a)}\ \widehat{N}_\star \simeq 61.5 + \ln\frac{\widehat{V}_{\rm HI}(\phi_\star)^{1/2}}{\widehat{V}_{\rm HI}(\phi_{\rm f})^{1/4}} + \frac{1}{2}f_R(\phi_\star) \quad\text{and}\quad \text{(b)}\ A_{\rm s}^{1/2} \simeq 4.627\cdot 10^{-5}\,, \tag{3.19}$$

where we consider in Eq. (3.19a) an equation-of-state parameter $w_{\rm int} = 1/3$ correspoding to quartic potential which is expected to approximate rather well $\widehat{V}_{\rm HI}$ for $\phi \ll 1$. We, thus, restrict λ and $\phi_\star$ and compute the model predictions via Eqs. (3.18a) and (3.18b) 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 ΛCDM+r model, i.e.,

$$\text{(a)}\ n_{\rm s} = 0.968 \pm 0.009 \quad\text{and}\quad \text{(b)}\ r \le 0.07, \tag{3.20}$$

at 95% *confidence level* (c.l.) with $|a_{\rm s}| \ll 0.01$.

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

$$\Phi \to \Phi/\sqrt{c_-},\ \ \bar{\Phi} \to \bar{\Phi}/\sqrt{c_-} \quad\text{and}\quad S \to S, \tag{3.21}$$

W in Eq. (2.2b) depends on λ/c_- and the K's in Eq. (2.3a) – (2.3c) depend on $r_\pm$. As a consequence, $\widehat{V}_{\rm HI}$ depends exclusively on λ/c_- and $r_\pm$. Since the λ/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 $n_{\rm s} - r_{0.002}$ plane, where $r_{0.002} =$

$16\widehat{\varepsilon}(\widehat{\phi}_{0.002})$ with $\widehat{\phi}_{0.002}$ being the value of $\widehat{\phi}$ when the scale $k = 0.002/\mathrm{Mpc}$, which undergoes $\widehat{N}_{0.002} = \widehat{N}_\star + 3.22$ e-foldings during HI, crosses the horizon of HI. We depict the theoretically allowed values with a solid line with the variation of $r_\pm$ shown along it. For low enough $r_\pm$'s – i.e. $r_\pm \leq 0.0005$ – the line reaches $(n_\mathrm{s}, r_{0.002}) \simeq (0.947, 0.28)$ obtained within *minimal* quartic inflation defined for $c_+ = 0$. Increasing $r_\pm$ the line enters the observationally allowed regions and terminates for $r_\pm \simeq 0.5$, beyond which Eq. (3.7) is violated. Along this line we find – consistently with the analytic formulas of Sec. 3.3

$$0.048 \lesssim \frac{r_\pm}{0.1} \lesssim 5,\;\; 9.64 \lesssim \frac{n_\mathrm{s}}{0.1} \lesssim 9.72,\;\; 0.7 \lesssim \frac{r}{0.01} \lesssim 8.1 \;\text{ and }\; 0.17 \lesssim 10^5 \frac{\lambda}{c_-} \lesssim 3.13\,. \tag{3.22}$$

Moreover $a_\mathrm{s} \simeq -(5-6)\cdot 10^{-4}$ 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 $r \gtrsim 0.007$. Had we employed $K = K_1$, the line in Fig. 1-(a) would have been shortened until $r_\pm \simeq 0.33$ yielding $r_{0.002} \gtrsim 0.0084$. 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:

$$r_\pm = 0.025 \;\text{ resulting to }\; (n_\mathrm{s}, r) = (0.969, 0.033)\,. \tag{3.23}$$

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

4. Higgs Inflation and μ Term of MSSM

A byproduct of our setting is the derivation of μ term of MSSM, as shown in Sec. 4.2, consistently with the low-energy phenomenology of MSSM – see Sec. 4.3. This construction is based on the SUSY potential found in Sec. 4.1. Hereafter we restore units, i.e., we take $m_\mathrm{P} = 2.433\cdot 10^{18}$ GeV.

4.1 SUSY Potential

The SUSY limit V_SUSY of $\widehat{V}_\mathrm{HI}$ in Eq. (3.4) is given by

$$V_\mathrm{SUSY} = \widetilde{K}^{\alpha\bar\beta} W_{\mathrm{HI}\alpha} W^*_{\mathrm{HI}\bar\beta} + \frac{g^2}{2}\Sigma_\mathrm{a} \mathrm{D}_\mathrm{a}\mathrm{D}_\mathrm{a}\,, \tag{4.1a}$$

where $\widetilde{K}$ is the limit of the K's in Eqs. (2.3a) – (2.3c) for $m_\mathrm{P} \to \infty$. Focusing on the $S - \bar\Phi - \Phi$ system we find

$$\widetilde{K} = c_- F_- - N c_+ F_+ + |S|^2. \tag{4.1b}$$

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

$$V_\mathrm{SUSY} = \lambda^2 \left|\bar\Phi\Phi - \frac{1}{4}M^2\right|^2 + \frac{\lambda^2}{c_-(1-Nr_\pm)} S^2\left(|\Phi|^2 + |\bar\Phi|^2\right) + \frac{g^2}{2} c_-^2 (1-Nr_\pm)^2 \left(|\Phi|^2 - |\bar\Phi|^2\right)^2.$$

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

$$\langle S\rangle = 0 \;\text{ and }\; |\langle\Phi\rangle| = |\langle\bar\Phi\rangle| = M/2\,. \tag{4.2}$$

As a consequence, $\langle\Phi\rangle$ and $\langle\bar\Phi\rangle$ 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 μ Term of MSSM

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

$$V_{\rm soft}=\left(\lambda A_\lambda S\bar\Phi\Phi+\lambda_\mu A_\mu SH_uH_d+\lambda_{iN^c}A_{iN^c}\Phi\widetilde N_i^{c2}-{\rm a}_S S\lambda M^2/4+{\rm h.c.}\right)+m_\gamma^2|X^\gamma|^2\,, \tag{4.3}$$

where $m_\gamma, A_\lambda, A_\mu, A_{iN^c}$ and ${\rm a}_S$ are soft SUSY breaking mass parameters. Rotating S in the real axis by an appropriate R-transformation, choosing conveniently the phases of A_λ and ${\rm a}_S$ so as the total low energy potential $V_{\rm tot}=V_{\rm SUSY}+V_{\rm soft}$ to be minimized – see Eq. (4.1c) – and substituting in $V_{\rm soft}$ the SUSY *vacuum expectation values* (v.e.vs) of Φ and $\bar\Phi$ from Eq. (4.2) we get

$$\langle V_{\rm tot}(S)\rangle=\lambda^2M^2S^2/2c_-(1-Nr_\pm)-\lambda{\rm a}_{3/2}m_{3/2}M^2S, \tag{4.4a}$$

where we take into account that $m_S\ll M$ and we set $|A_\lambda|+|{\rm a}_S|=2{\rm a}_{3/2}m_{3/2}$ with $m_{3/2}$ being the $\widetilde G$ mass and ${\rm a}_{3/2}>0$ a parameter of order unity which parameterizes our ignorance for the dependence of $|A_\lambda|$ and $|{\rm a}_S|$ 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

$$\frac{d}{dS}\langle V_{\rm tot}(S)\rangle=0\ \Rightarrow\ \langle S\rangle\simeq{\rm a}_{3/2}m_{3/2}c_-(1-Nr_\pm)/\lambda. \tag{4.4b}$$

The generated μ term from the third term in the r.h.s of Eq. (2.2b) is

$$\mu=\lambda_\mu\langle S\rangle\simeq\lambda_\mu{\rm a}_{3/2}m_{3/2}c_-(1-Nr_\pm)/\lambda\,. \tag{4.5}$$

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

4.3 Connection with the MSSM Phenomenology

The SUSY breaking effects, considered in Eq. (4.3), explicitly break $U(1)_R$ to a subgroup, $\mathbb{Z}_2^R$ 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 $\mathbb{Z}_2^R$ remains unbroken, since $\langle S\rangle$ in Eq. (4.4b) also breaks spontaneously $U(1)_R$ to $\mathbb{Z}_2^R$ and so no disastrous domain walls are formed. Combining $\mathbb{Z}_2^R$ with the $\mathbb{Z}_2^{\rm f}$ 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

$${\rm sign}\mu,\ \ \tan\beta=\langle H_u\rangle/\langle H_d\rangle,\ \ M_{1/2},\ \ m_0\ \text{ and }\ A_0, \tag{4.6}$$

where signμ is the sign of μ, and the three last mass parameters denote the common gaugino mass, scalar mass and trilinear coupling constant, respectively, defined (normally) at $M_{\rm GUT}$. The parameter

CMSSM	$\lvert A_0\rvert$	m_0	$\lvert\mu\rvert$	$a_{3/2}$	$\lambda_\mu\ (10^{-6})$	
REGION	(TeV)	(TeV)	(TeV)		$K=K_1$	$K=K_2,\ K_3$
A/H Funnel	9.9244	9.136	1.409	1.086	0.6223	0.607
$\tilde{\tau}_1-\chi$ Coannihilation	1.2271	1.476	2.62	0.831	*9.36*	9.12
$\tilde{t}_1-\chi$ Coannihilation	9.965	4.269	4.073	2.33	1.794	1.75
$\tilde{\chi}_1^{\pm}-\chi$ Coannihilation	9.2061	9.000	0.983	1.023	0.468	0.456

Table 3: *The required λ_μ values which render our models compatible with the best-fit points in the CMSSM, as found in Ref. [8], for the assumptions of Eq. (4.7) $K=K_1$ or $K=K_2$ and K_3 with $N_X=2$ and $r_\pm=0.025$.*

$|\mu|$ 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 (χ) – $\tilde{\tau}_1,\tilde{t}_1$ and $\tilde{\chi}_1^{\pm}$ 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

$$m_0=m_{3/2}\ \text{ and }\ |A_0|=|A_\lambda|=|\mathrm{a}_S|, \tag{4.7}$$

we can derive first $\mathrm{a}_{3/2}$ and then the λ_μ values which yield the phenomenologically desired $|\mu|$ – ignoring renormalization group effects. The outputs of our computation are 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 λ_μ 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

$$\lambda_\mu\lesssim 6.6\cdot 10^{-6}\ \text{ for }\ K=K_1\ \text{ and }\ \lambda_\mu\lesssim 1.1\cdot 10^{-5}\ \text{ for }\ K=K_2\ \text{ and }\ K_3\,. \tag{4.8}$$

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 $T_{\rm rh}$ – see Sec. 5.2. Indeed, $m_{3/2}$ as low as 1 TeV becomes cosmologically safe, under the assumption of the unstable $\widetilde{G}$, for the $T_{\rm rh}$ 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 $\widetilde{G}$ 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 $\widehat{V}_{\rm HI}$. The (canonically normalized) inflaton,

$$\widehat{\delta\phi} = \langle J\rangle\delta\phi \;\text{ with }\; \delta\phi = \phi - M \;\text{ and }\; \langle J\rangle = \sqrt{\langle\kappa_+\rangle} = \sqrt{c_-(1-Nr_\pm)} \tag{5.1}$$

acquires mass, at the SUSY vacuum in Eq. (4.2), which is given by

$$\widehat{m}_{\delta\phi} = \left\langle \widehat{V}_{{\rm HI},\widehat{\phi}\widehat{\phi}}\right\rangle^{1/2} = \left\langle \widehat{V}_{{\rm HI},\phi\phi}/J^2\right\rangle^{1/2} \simeq \frac{\lambda M}{\sqrt{2c_-\,(1-Nr_\pm)}}\,, \tag{5.2}$$

where the last (approximate) equality above is valid only for $r_\pm \ll 1/N$ – see Eqs. (3.6) and (3.8). As we see, $\widehat{m}_{\delta\phi}$ 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 $\langle\bar\Phi\rangle$ in Eq. (4.2) take the values dictated by the unification of the MSSM gauge coupling constants, despite the fact that $U(1)_{B-L}$ gauge symmetry does not disturb this unification and M could be much lower. In particular, the unification scale $M_{\rm GUT} \simeq 2\cdot 10^{16}$ GeV can be identified with M_{BL} – see Table 2 – at the SUSY vacuum in Eq. (4.2), i.e.,

$$\frac{\sqrt{c_-(\langle f_R\rangle - Nr_\pm)}gM}{\sqrt{\langle f_R\rangle}} = M_{\rm GUT} \;\Rightarrow\; M \simeq M_{\rm GUT}/g\sqrt{c_-\,(1-Nr_\pm)} \tag{5.3}$$

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 $\widehat{m}_{\delta\phi}$ remains constant for fixed $r_\pm$ since λ/c_- is fixed too – see Eq. (3.17). Particularly, along the line in Fig. 1-(a) we obtain

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

$$3.46\cdot 10^{10} \lesssim \widehat{m}_{\delta\phi}/{\rm GeV} \lesssim 4.2\cdot 10^{13} \quad \text{for }\; K = K_2 \;\text{ and }\; K_3\,. \tag{5.4b}$$

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

$$T_{\rm rh} = \left(72/5\pi^2 g_*\right)^{1/4}\left(\widehat{\Gamma}_{\delta\phi}m_{\rm P}\right)^{1/2} \;\text{ with }\; \widehat{\Gamma}_{\delta\phi} = \widehat{\Gamma}_{\delta\phi\to N^c_iN^c_i} + \widehat{\Gamma}_{\delta\phi\to H_uH_d}\,. \tag{5.5}$$

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

$$\widehat{\Gamma}_{\delta\phi\to N^c_iN^c_i} = \frac{g^2_{iN^c}}{16\pi}\widehat{m}_{\delta\phi}\left(1-\frac{4M^2_{iN^c}}{\widehat{m}^2_{\delta\phi}}\right)^{3/2} \;\text{ with }\; g_{iN^c} = \frac{\lambda_{iN^c}}{\langle J\rangle}\left(1-3c_+\frac{N}{2}\frac{M^2}{m^2_{\rm P}}\right), \tag{5.6a}$$

$$\widehat{\Gamma}_{\delta\phi\to H_uH_d} = \frac{2}{8\pi}g^2_H\widehat{m}_{\delta\phi} \;\text{ with }\; g_H = \frac{\lambda_\mu}{\sqrt{2}}\left(1-2c_+\frac{M^2}{m^2_{\rm P}}\right) \tag{5.6b}$$

arising from the lagrangian terms

$$\mathscr{L}_{\widehat{\delta\phi}\to N^c_iN^c_i} = -\frac{1}{2}e^{K/2m^2_{\rm P}}W_{{\rm HI},N^c_iN^c_i}N^c_iN^c_i \;+\; {\rm h.c.} = g_{iN^c}\widehat{\delta\phi}\;\left(N^c_iN^c_i \;+\; {\rm h.c.}\right)+\cdots, \tag{5.6c}$$

$$\mathscr{L}_{\widehat{\delta\phi}\to H_uH_d} = -e^{K/m^2_{\rm P}}K^{SS^*}\left|W_{{\rm HI},S}\right|^2 = -g_H\widehat{m}_{\delta\phi}\widehat{\delta\phi}\left(H^*_uH^*_d \;+\; {\rm h.c.}\right)+\cdots \tag{5.6d}$$

describing $\widehat{\delta\phi}$ decay into a pair of N^c_j with masses $M_{jN^c} = \lambda_{jN^c}M$ and H_u and H_d respectively. Note that the decay modes into three MSSM (s)-particles through a typical trilinear term in Eq. (2.2a) is suppressed, since they arise from non-renormalizable interactions proportional to $M/m_{\rm P} \ll 1$ [1].

5.2 Lepton-Number and Gravitino Abundances

For $T_{\rm rh} < M_{iN^c}$, the out-of-equilibrium decay of N^c_i generates a lepton-number asymmetry (per N^c_i decay), ε_i. The resulting lepton-number asymmetry is partially converted through sphaleron effects into a yield of the observed BAU

$$Y_B = -0.35\cdot\frac{5}{2}\frac{T_{\rm rh}}{\widehat{m}_{\delta\phi}}\sum_i\frac{\widehat{\Gamma}_{\delta\phi\to N^c_iN^c_i}}{\widehat{\Gamma}_{\delta\phi}}\varepsilon_i \ \text{ with } \ \varepsilon_i=\sum_{j\neq i}\frac{\mathsf{Im}\left[(m^\dagger_{\rm D}m_{\rm D})^2_{ij}\right]}{8\pi\langle H_u\rangle^2(m^\dagger_{\rm D}m_{\rm D})_{ii}}\Big(F_{\rm S}\left(x_{ij},y_i,y_j\right)+F_{\rm V}(x_{ij})\Big), \quad (5.7)$$

where $\langle H_u\rangle \simeq 174$ GeV, for large $\tan\beta$, $m_{\rm D}$ is the Dirac mass matrix of neutrinos ,ν_i, and $F_{\rm S}$ [$F_{\rm V}$] 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]

$$Y_B = \left(8.64^{+0.15}_{-0.16}\right)\cdot 10^{-11}. \quad (5.8)$$

The validity of Eq. (5.7) requires that the $\widehat{\delta\phi}$ decay into a pair of N^c_i's is kinematically allowed for at least one species of the N^c_i's and also that there is no erasure of the produced Y_L due to N^c_1 mediated inverse decays and $\Delta L = 1$ scatterings. These prerequisites are ensured if we impose

$$(\mathsf{a})\ \widehat{m}_{\delta\phi} \geq 2M_{1N^c} \ \text{ and } \ (\mathsf{b})\ M_{1N^c} \gtrsim 10T_{\rm rh}. \quad (5.9)$$

The quantity ε_i can be expressed in terms of the Dirac masses of ν_i, $m_{i\rm D}$, arising from the third term of Eq. (2.2b). Employing the seesaw formula we can then obtain the light-neutrino mass matrix m_ν in terms of $m_{i\rm D}$ and M_{iN^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, $\Delta m^2_{21} = 7.56\cdot 10^{-5}$ eV^2 and $\Delta m^2_{31} = 2.55\cdot 10^{-3}$ eV^2 $\left[\Delta m^2_{31} = 2.49\cdot 10^{-3}\text{ eV}^2\right]$, on the mixing angles, $\sin^2\theta_{12} = 0.321$, $\sin^2\theta_{13} = 0.02155$ $\left[\sin^2\theta_{13} = 0.0214\right]$ and $\sin^2\theta_{23} = 0.43$ $\left[\sin^2\theta_{23} = 0.596\right]$ and the CP-violating Dirac phase $\delta = 1.4\pi$ $[\delta = 1.44\pi]$ for *normal [inverted] ordered* (NO [IO]) *neutrino masses*, $m_{i\nu}$'s. Furthermore, the sum of $m_{i\nu}$'s is bounded from above at 95% c.l. by the data [6]

$$\textstyle\sum_i m_{i\nu} \leq 0.23 \text{ eV}. \quad (5.10)$$

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

$$Y_{3/2} \simeq 1.9\cdot 10^{-22}\ T_{\rm rh}/\text{GeV}, \quad (5.11)$$

where we take into account only thermal production of $\widetilde{G}$, and assume that $\widetilde{G}$ 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 $\widetilde{G}$ decays with a tiny hadronic branching ratio, we have

$$Y_{3/2} \lesssim \begin{cases}10^{-14}\\ 10^{-13}\\ 10^{-12}\end{cases} \text{for } m_{3/2} \simeq \begin{cases}0.69\text{ TeV}\\ 10.6\text{ TeV}\\ 13.5\text{ TeV}\end{cases} \text{implying } T_{\rm rh} \lesssim 5.3\cdot\begin{cases}10^7\text{ GeV},\\ 10^8\text{ GeV},\\ 10^9\text{ GeV}.\end{cases} \quad (5.12)$$

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

PARAMETERS	CASES						
	A	B	C	D	E	F	G
	NORMAL HIERARCHY		ALMOST DEGENERACY			INVERTED HIERARCHY	
LOW SCALE PARAMETERS							
$m_{1\nu}/0.1$ eV	0.05	0.1	0.5	0.7	0.7	0.5	0.49
$m_{2\nu}/0.1$ eV	0.1	0.13	0.51	0.7	0.7	0.51	0.5
$m_{3\nu}/0.1$ eV	0.5	0.51	0.7	0.86	0.5	0.1	0.05
$\sum_i m_{i\nu}/0.1$ eV	0.65	0.74	1.7	2.3	1.9	1.1	1
φ_1	$-\pi/5$	$-\pi/2$	π	$\pi/9$	0	$-3\pi/4$	$\pi/2$
φ_2	π	0	$\pi/3$	π	$\pi/2$	$5\pi/4$	$-\pi/2$
LEPTOGENESIS-SCALE PARAMETERS							
$m_{1\mathrm{D}}/0.1$ GeV	2	5	10	10	1.3	5	6
$m_{2\mathrm{D}}/$GeV	6	1.97	3.9	10	9	0.715	1.1
$m_{3\mathrm{D}}/$GeV	100	150	170	168	202	100	199
$M_{1N^c}/10^{10}$ GeV	1.0	3.3	2.85	3.3	2.98	0.45	1.23
$M_{2N^c}/10^{10}$ GeV	6.9	13.6	26.5	111.4	13.9	2.76	2.8
$M_{3N^c}/10^{14}$ GeV	2.9	4.9	2.2	1.2	3.7	2.7	27.2
OPEN DECAY CHANNELS OF THE INFLATON, $\widehat{\delta\phi}$, INTO N_i^c							
$\widehat{\delta\phi}\ \to$	N_1^c	N_1^c	N_1^c	N_1^c	N_1^c	$N_{1,2}^c$	$N_{1,2}^c$
$\widehat{\Gamma}_{\delta\phi\to N_i^cN_i^c}/\widehat{\Gamma}_{\delta\phi}$ (%)	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 $T_{\rm rh}$ AND $\widetilde{G}$-YIELD							
$T_{\rm rh}/10^7$ 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 $N_X=2$, $r_\pm=0.025$ and $\lambda_\mu=10^{-6}$.*

5.3 Results

Confronting with observations Y_B and $Y_{3/2}$ which depend on $\widehat{m}_{\delta\phi}$, $T_{\rm rh}$, M_{iN^c} and $m_{i\mathrm{D}}$'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_{iN^c}'s by using as inputs

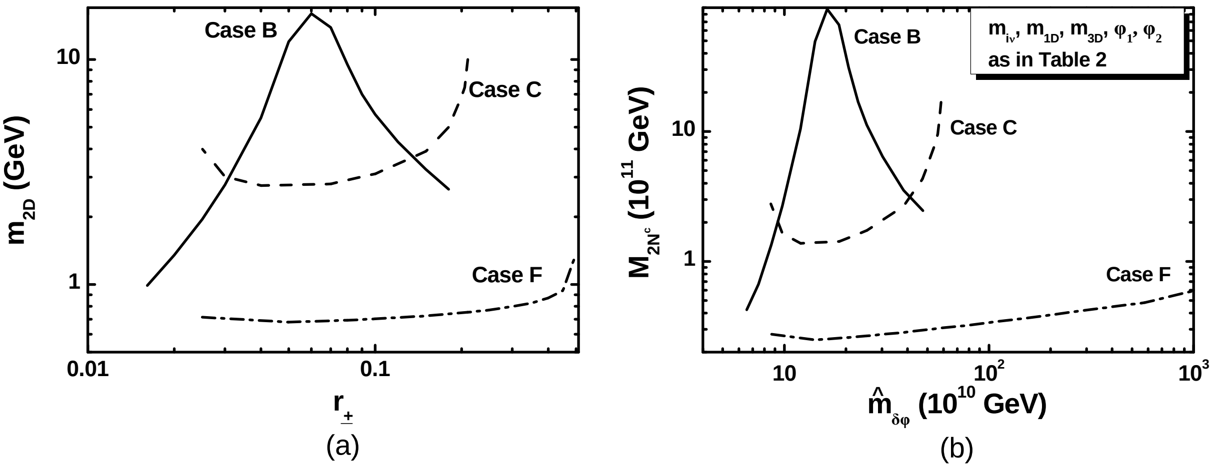

Figure 2: *Contours, yielding the central Y_B in Eq. (5.8) consistently with the inflationary requirements, in the (a) $r_{\pm} - m_{2\mathrm{D}}$ and (b) $\widehat{m}_{\delta\phi} - M_{2N^c}$ plane. We take $K = K_2$ or K_3 with $N_X = 2$, $\lambda_\mu = 10^{-6}$ and the values of $m_{i\nu}$, $m_{1\mathrm{D}}$, $m_{3\mathrm{D}}$, φ_1 and φ_2 which correspond to the cases B (solid line), C (dashed line) and F (dot-dashed line) of Table 4.*

the $m_{i\mathrm{D}}$'s, a reference mass of the ν_i's – $m_{1\nu}$ for NO $m_{i\nu}$'s, or $m_{3\nu}$ for IO $m_{i\nu}$'s –, the two Majorana phases φ_1 and φ_2 of the PMNS matrix, and the best-fit values for the low energy parameters of neutrino physics mentioned in Sec. 5.2. In our numerical code, we also estimate [1] the RG evolved values of the latter parameters at the scale of nTL, $\Lambda_L = \widehat{m}_{\delta\phi}$, by considering the MSSM with $\tan\beta \simeq 50$ as an effective theory between Λ_L and the soft SUSY breaking scale, $M_{\mathrm{SUSY}} = 1.5$ TeV. We evaluate the M_{iN^c}'s at Λ_L, and we neglect any possible running of the $m_{i\mathrm{D}}$'s and M_{iN^c}'s. The so obtained M_{iN^c}'s clearly correspond to the scale Λ_L.

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. (3.23) and $\lambda_\mu = 10^{-6}$ in accordance with Eqs. (3.10a) and (3.10b) with $N_X = 2$. We obtain $\widehat{m}_{\delta\phi} = 8.9 \cdot 10^{10}$ GeV for $K = K_1$ and $\widehat{m}_{\delta\phi} = 8.6 \cdot 10^{10}$ 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) $m_{i\nu}$'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 $m_{i\mathrm{D}}$'s are free parameters. This fact facilitates the fulfilment of Eq. (5.9b) since $m_{1\mathrm{D}}$ affects heavily M_{1N^c}. Care is also taken so that the perturbativity of λ_{iN^c} holds, i.e., $\lambda_{iN^c}^2/4\pi \leq 1$. The inflaton $\widehat{\delta\phi}$ decays mostly into N_1^c's – see cases A – E. In all cases $\widehat{\Gamma}_{\delta\phi\to N_i^c N_i^c} < \widehat{\Gamma}_{\delta\phi\to H_u H_d}$ and so the ratios $\widehat{\Gamma}_{\delta\phi\to N_i^c N_i^c}/\widehat{\Gamma}_{\delta\phi}$ introduce a considerable reduction in the derivation of Y_B. In Table 4 we also display the values of T_{rh}, the majority of which are close to $3 \cdot 10^7$ GeV, and the corresponding $Y_{3/2}$'s, which are consistent with Eq. (5.12) for $m_{3/2} \gtrsim 1$ TeV. These values are in nice agreement with the ones needed for the solution of the μ problem of MSSM – see, e.g., Table 3.

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 $\widehat{m}_{\delta\phi}$ and adjusting conveniently $m_{2\mathrm{D}}$ or M_{2N^c} – 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. 2-(b). In particular, we use solid, dashed and dot-dashed line when the remaining inputs – i.e. $m_{i\nu}$, $m_{1\mathrm{D}}$, $m_{3\mathrm{D}}$, φ_1, and φ_2 – correspond to the cases B, C and F of Table 4, 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 $m_{2\mathrm{D}}$ 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

$$4.4 \lesssim Y_{\widetilde{G}}/10^{-15} \lesssim 228 \;\text{ and }\; 0.23 \lesssim T_{\mathrm{rh}}/10^8\mathrm{GeV} \lesssim 12 \;\text{ with }\; 6.5 \lesssim \widehat{m}_{\delta\phi}/10^{10} \lesssim 4241\,. \tag{5.13}$$

As a bottom line, nTL not only is a realistic possibility within our setting but also it can be comfortably reconciled with the $\widetilde{G}$ constraint even for $m_{3/2} \sim 1$ 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 μ problem of MSSM, provided that λ_μ is somehow small; (iv) they allow for baryogenesis via nTL compatible with $\widetilde{G}$ constraints and neutrino data. In particular, we may have $m_{3/2} \sim 1$ TeV, with the inflaton decaying mainly to N_1^c and N_2^c – we obtain M_{iN^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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