



Supersymmetric $B - L$ inflation near the conformal coupling



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ABSTRACT

We investigate a novel scenario of cosmological inflation in a gauged $B - L$ extended minimal supersymmetric Standard Model with R-symmetry. We use a noncanonical Kähler potential and a superpotential, both preserving the R-symmetry to construct a model of slow-roll inflation. The model is controlled by two real parameters: the nonminimal coupling ξ that originates from the Kähler potential, and the breaking scale v of the $U(1)_{B-L}$ symmetry. We compute the spectrum of the cosmological microwave background radiation and show that the prediction of the model fits well the recent Planck satellite observation for a wide range of the parameter space. We also find that the typical reheating temperature of the model is low enough to avoid the gravitino problem but nevertheless allows sufficient production of the baryon asymmetry if we take into account the effect of resonance enhancement. The model is free from cosmic strings that impose stringent constraints on generic $U(1)_{B-L}$ based scenarios, as in our scenario the $U(1)_{B-L}$ symmetry is broken from the onset.

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1. Introduction

Recent observations of the cosmological microwave background (CMB) [1–4] impose stringent restrictions on models of inflation. For example, the minimally coupled $m^2\phi^2$ and $\lambda\phi^4$ chaotic models that have served as simple benchmark models for decades are now in tension. Inflationary models with nonminimal coupling $\xi\phi^2 R$, where ϕ is a scalar field (inflaton) and R the scalar curvature, are less constrained, and in fact the predictions of some such models have been shown to fit extremely well the current data [5]. It is well known that a nonminimally coupled model (in the Jordan frame) can be Weyl-rescaled to a minimally coupled model (the Einstein frame), and hence it is meaningful to discuss the former only when the original Lagrangian has significance in its own right, e.g. when the Lagrangian is that of a particle physics model such as the Standard Model (SM). The nonminimally coupled SM Higgs inflation [6–10] with the Jordan frame Higgs potential

$$V_J = \lambda(\phi^2 - v^2)^2, \quad (1)$$

provides just such an example. While the hierarchy problem inherited from the SM and a large dimensionless coupling $\xi \sim 10^4$

required for the consistency loom large and stay as a matter of debate [11–15], the simplicity and the observational viability are very attractive features. The hierarchy problem is known to be mitigated in a supersymmetric setup. Supersymmetric extensions of the Higgs inflation have been proposed in the next-to-minimal supersymmetric SM [16–18] and in the supersymmetric grand unified theory (GUT) [19,20] (see also [21]). As a closely related model, it is shown in [22,23] that the Higgs-lepton flat direction in the supersymmetric seesaw Lagrangian can realise observationally viable and phenomenologically consistent slow roll inflation with small nonminimal coupling $\xi \lesssim \mathcal{O}(1)$.

These popular Higgs inflation models employ positive¹ nonminimal coupling ξ . Interestingly, it is known that the same Jordan frame potential (1) with *negative* nonminimal coupling ξ also provides a model of inflation [24,25], in which an observationally viable case correspond to small field values $|\phi| \lesssim v$. The potential in the Einstein frame is

$$V_E \sim \left(\frac{\phi^2 - v^2}{1 - |\xi|\phi^2} \right)^2, \quad (2)$$

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¹ Our convention is such that the conformal coupling in 4 dimensions corresponds to $\xi = -\frac{1}{6}$.

Table 1The charges of the superfields under the symmetries of the SM gauge group, $U(1)_{B-L}$, and $U(1)_R$.

	Q	u^c	d^c	L	e^c	H_u	H_d	N^c	S	Φ_{\pm}
$SU(3)_c \times SU(2)_L \times U(1)_Y$	$(\mathbf{3}, \mathbf{2}, +\frac{1}{6})$	$(\bar{\mathbf{3}}, \mathbf{1}, -\frac{2}{3})$	$(\bar{\mathbf{3}}, \mathbf{1}, +\frac{1}{3})$	$(\mathbf{1}, \mathbf{2}, -\frac{1}{2})$	$(\mathbf{1}, \mathbf{1}, +1)$	$(\mathbf{1}, \mathbf{2}, +\frac{1}{2})$	$(\mathbf{1}, \mathbf{2}, -\frac{1}{2})$	$(\mathbf{1}, \mathbf{1}, 0)$	$(\mathbf{1}, \mathbf{1}, 0)$	$(\mathbf{1}, \mathbf{1}, 0)$
$U(1)_{B-L}$	$+\frac{1}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$	-1	$+1$	0	0	$+1$	0	± 2
$U(1)_R$	$+\frac{1}{2}$	$+\frac{1}{2}$	$+\frac{1}{2}$	0	$+1$	$+1$	$+1$	$+1$	$+2$	0

where the reduced Planck mass $M_P = (8\pi G)^{-1/2} = 2.44 \times 10^{18}$ GeV has been set to unity. The potential exhibits global minima at $\phi = \pm v$ and singularities at $\phi = \pm 1/\sqrt{|\xi|}$. Hence successful termination of inflation (exit from the slow roll at a global minimum) requires $v^2|\xi| < 1$. It would be interesting to know in which context of particle physics such a model of inflation may be implemented, and in particular, what the broken symmetry associated with the potential (2) can be. Certainly, ϕ cannot be the SM Higgs field as the value of v required for inflation is much larger than the electroweak scale. A natural guess might be that ϕ is a Higgs field responsible for breaking some extra gauge symmetry. It would be then important to examine whether the resulting cosmological scenario is phenomenologically consistent.

In this paper we point out the possibility that ϕ can be the Higgs field associated with the $U(1)_{B-L}$ symmetry which is spontaneously broken at an ultra high energy scale in the early Universe. The $U(1)_{B-L}$ symmetry is one of the global symmetries of the minimal supersymmetric Standard Model (MSSM), under which the quark, lepton, and Higgs superfields are charged by $+\frac{1}{3}$, -1 , and 0 units (see Table 1). We assume that the $U(1)_{B-L}$ gauge symmetry is spontaneously broken at a super-Planckian scale; such an assumption is acceptable on phenomenological grounds as the breaking scale is experimentally unconstrained except the rather mild LEP bound $\gtrsim 3$ TeV [26,27]. An important consequence of having the $U(1)_{B-L}$ symmetry is that three generations of right-handed neutrinos are necessary for anomaly cancellation. Thus our model necessarily involves the neutrino sector; this allows us to discuss the neutrino masses (via the seesaw mechanism [28]) and the baryon asymmetry of the Universe (via thermal [29] or nonthermal [30] leptogenesis) within the same model. Indeed, the $U(1)_{B-L}$ -extended SM is one of the leading candidates of the particle physics beyond the SM and there are well known inflationary models based on it [31–40] (see also [41,42]). The novelty of our scenario, in comparison to the existing ones, is simplicity of the construction and robustness of the prediction. Also, our model is free from overproduction of cosmic strings that generally afflicts the $U(1)_{B-L}$ -based inflationary models; in our scenario the $U(1)_{B-L}$ symmetry is already broken at the onset of inflation and there is no danger of producing topological defects during and after inflation. We shall consider supergravity-embedding as the nonminimal coupling of the inflaton naturally arises in such a framework.

The rest of the paper is organised as follows. In the next section we construct the model from the supergravity setup, and in Section 3 we discuss the inflationary dynamics. We compare the prediction of this model with the results of the Planck satellite observation in Section 4, and the post-inflationary physics is discussed in Section 5. We conclude the paper with comments in Section 6.

2. The model

The starting point of our model is the superpotential

$$W_{\text{eff}} = \kappa S(\Phi_+ \Phi_- - v^2), \quad (3)$$

where the superfields S , Φ_+ , Φ_- are singlets in the SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$, and carry 0 , $+2$, -2 units

of $U(1)_{B-L}$ charges. We choose $\kappa > 0$, $v > 0$ by field redefinition. The local $U(1)_{B-L}$ symmetry is broken by the vacuum expectation value v of the Φ_{\pm} fields. The model may be considered as a part of a supersymmetric SM whose superpotential is (for example)

$$W = \mu H_u H_d + y_u^{ij} u_i^c Q_j H_u + y_d^{ij} d_i^c Q_j H_d + y_e^{ij} e_i^c L_j H_d + \kappa S(\Phi_+ \Phi_- - v^2) + y_D^{ij} N_i^c L_j H_u + \lambda^{ij} \Phi_- N_i^c N_j^c. \quad (4)$$

The first line represents the MSSM and the last two terms are responsible for the seesaw mechanism and leptogenesis. Here, Q , u^c , d^c , L , e^c , H_u , H_d are the MSSM superfields, N^c the right-handed neutrino superfields, y 's are the Yukawa couplings and μ is the MSSM μ parameter. The family indices are $i, j = 1, 2, 3$. There are three right-handed neutrinos necessary for anomaly cancellation. The Majorana Yukawa coupling λ^{ij} controls the seesaw scale; we shall discuss it in a later section. With $+2, 0, 0$ units of R-charges assigned to the S , Φ_+ , Φ_- fields, the superpotential is also invariant under the $U(1)_R$ symmetry.² In Table 1 we list the SM gauge group, $U(1)_{B-L}$ and $U(1)_R$ charges assigned to the superfields appearing in the superpotential (4). Note that (3) is the most general renormalisable superpotential for S and Φ_{\pm} that is compatible with these symmetries. For supergravity embedding in the superconformal framework, we shall use a slightly noncanonical Kähler potential $K = -3\Phi$, where

$$\Phi = 1 - \frac{1}{3}(|\Phi_+|^2 + |\Phi_-|^2 + |S|^2) + \frac{\gamma}{2}(\Phi_+ \Phi_- + \Phi_+^* \Phi_-^*) + \frac{\zeta}{3}|S|^4. \quad (5)$$

This preserves the $U(1)_R$ symmetry and contains two real parameters γ and ζ .

The scalar potential is found by the standard supergravity computation [44]. We take the D-flatness direction $|\Phi_+| = |\Phi_-| \equiv \frac{1}{2}\varphi$ and define

$$\Phi_+ = \frac{1}{2}\varphi e^{i\theta_+}, \quad \Phi_- = \frac{1}{2}\varphi e^{i\theta_-}, \quad (6)$$

with real scalar fields φ , θ_{\pm} . The F-term scalar potential is then³

$$V_F = \frac{1}{2}\kappa^2 |S|^2 \varphi^2 + \frac{\kappa^2}{1-4\zeta|S|^2} \left(\frac{\varphi^4}{16} - \frac{1}{2}\varphi^2 v^2 \cos\theta + v^4 \right) - \kappa^2 |S|^2 \frac{|\left(\frac{3}{4}\gamma - \frac{\zeta|S|^2 e^{i\theta}}{1-4\zeta|S|^2}\right)\frac{\varphi^2}{2} - 2\left(1 - \frac{\zeta|S|^2}{1-4\zeta|S|^2}\right)v^2|^2}{3 - \frac{3\gamma}{8}\varphi^2 \cos\theta + \frac{9\gamma^2}{32}\varphi^2 + \frac{\zeta|S|^4}{1-4\zeta|S|^2}}, \quad (7)$$

where $\theta \equiv \theta_+ + \theta_-$. Note that V_F is invariant under the phase shift $S \rightarrow e^{i\alpha} S$, reflecting the unbroken $U(1)_R$ symmetry.

Let us comment on the F-term hybrid inflation (FHI) models [31–38] which share some similarity with ours. The FHI models

² These R-charges are what is called $R' = R - \frac{1}{2}L$ in [43]. In our model both R and R' are conserved.

³ The super-Planckian inflaton values imply that higher dimensional operators are not negligible. To avoid deformation of the potential due to such operators, some degree of fine-tuning is unavoidable.

are based on the same superpotential (3) representing the spontaneously broken local $U(1)_{B-L}$ symmetry, but usually the canonical Kähler potential is assumed.⁴ In these models the major role is played by the S field whereas the role played by φ (called a waterfall field) is minor. At the tree level the scalar spectral index n_s of the FHI models is typically enhanced (blue). The slightly red n_s compatible with the current observations can be obtained by including radiative and supergravity correction terms (see [51,52] for up-to-date accounts). Here in our model we take a different trajectory from the FHI models. It can be shown that for large enough ζ , the field S is stabilised at $S=0$ and its dynamics can be neglected (cf. [17,18]). Then the potential (7) and Φ of (5) simplify to

$$V_F = \frac{\kappa^2}{16} (\varphi^4 - 8\varphi^2 v^2 \cos\theta + 16v^4), \quad (8)$$

$$\Phi = 1 + \left(\frac{\gamma}{4} \cos\theta - \frac{1}{6} \right) \varphi^2. \quad (9)$$

Examining the F-term scalar potential in the Einstein frame $V_E = \Phi^{-2} V_F$, it can be checked that the phase is stable at $\theta=0$ (we will be concerned with the parameter region $\gamma > 0$; see below). Thus one may further ignore the dynamics of θ . The system then reduces to a single field model for which the scalar-gravity part of the Lagrangian is written as

$$\mathcal{L}_J = \sqrt{-g_J} \left[\frac{1}{2} \Phi R_J - \frac{1}{2} g_J^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi - V_J \right], \quad (10)$$

where the subscript J stands for the Jordan frame and

$$\Phi = 1 + \xi \varphi^2, \quad (11)$$

$$V_J = V_F = \kappa^2 \left(\frac{\varphi^2}{4} - v^2 \right)^2, \quad (12)$$

$$\xi = \frac{\gamma}{4} - \frac{1}{6}. \quad (13)$$

The Lagrangian in the Einstein frame is related to the one in the Jordan frame by the Weyl transformation $g_{\mu\nu}^E = \Phi g_{\mu\nu}^J$ and is written as

$$\mathcal{L}_E = \sqrt{-g_E} \left[\frac{1}{2} R_E - \frac{1}{2} g_E^{\mu\nu} \partial_\mu \hat{\varphi} \partial_\nu \hat{\varphi} - V_E \right], \quad (14)$$

where R_E is the scalar curvature in the Einstein frame, $g_E^{\mu\nu} = (g_{\mu\nu}^E)^{-1}$, $V_E = \Phi^{-2} V_J$, and $\hat{\varphi}$ is the canonically normalised scalar fields in the Einstein frame which is related to φ via

$$d\hat{\varphi} = \frac{\sqrt{1 + \xi\varphi^2 + 6\xi^2\varphi^2}}{1 + \xi\varphi^2} d\varphi. \quad (15)$$

The slow roll parameters defined in the Einstein frame are

$$\epsilon_V = \frac{1}{2} \left(\frac{1}{V_E} \frac{dV_E}{d\hat{\varphi}} \right)^2, \quad \eta_V = \frac{1}{V_E} \frac{d^2 V_E}{d\hat{\varphi}^2}, \quad (16)$$

$$\xi_V^2 = \frac{1}{V_E^2} \frac{dV_E}{d\hat{\varphi}} \frac{d^3 V_E}{d\hat{\varphi}^3}.$$

Using the original field φ these are expressed as

$$\epsilon_V = \frac{1}{2} \left(\frac{V_E'}{V_E \hat{\varphi}'} \right)^2, \quad \eta_V = \frac{V_E''}{V_E (\hat{\varphi}')^2} - \frac{V_E' \hat{\varphi}''}{V_E (\hat{\varphi}')^3}, \quad (17)$$

$$\xi_V^2 = \frac{V_E'}{V_E^2} \left(3 \frac{V_E' (\hat{\varphi}'')^2}{(\hat{\varphi}')^6} - \frac{V_E' \hat{\varphi}'''}{(\hat{\varphi}')^5} - 3 \frac{V_E'' \hat{\varphi}''}{(\hat{\varphi}')^5} + \frac{V_E'''}{(\hat{\varphi}')^4} \right),$$

where the prime means $' \equiv d/d\varphi$.

3. Inflationary dynamics

In the last section we obtained the single field inflationary model from the gauged $B-L$ extended MSSM. The inflaton potential in the Einstein frame is

$$V_E = \frac{\kappa^2}{16} \left(\frac{\varphi^2 - 4v^2}{1 + \xi\varphi^2} \right)^2, \quad (18)$$

which is essentially the one (2) discussed in the introduction. The potential (18) has supersymmetric vacua at $\varphi = \pm 2v$ and singularities at $\varphi = \pm 1/\sqrt{|\xi|}$. Without losing generality we shall focus on positive v and positive φ . In this paper we will be interested in the inflationary scenario with negative ξ [24,25]. The initial value of the inflaton is between $0 < \varphi < 2v$, namely, between the local maximum of the potential $V_E(0) = \kappa^2 v^4$ and the supersymmetric vacuum at $\varphi = 2v$. This is analogous to the *new inflation* model, or even newer, hilltop type models [53]. For successful termination of the slow roll the supersymmetric vacuum $\varphi = 2v$ should not be hidden behind a singularity; this requires $2v < 1/\sqrt{|\xi|}$. The physics beyond the singularity is of no interest to us, as it is the antigravity regime where the Newton constant becomes negative [54,55].

In our model there are three tunable parameters κ , v and ξ . These are constrained by the amplitude of the density perturbation of the comoving CMB scale. For definiteness we use the maximum likelihood value $A_s(k_0) = 2.215 \times 10^{-9}$ from the Planck satellite observation [3] with the pivot scale at $k_0 = 0.05 \text{ Mpc}^{-1}$. With this $A_s(k)$ the power spectrum of the curvature perturbation $\mathcal{P}_R = V_E/24\pi^2 \epsilon_V$ at the horizon exit of the comoving scale is normalised as $A_s(k) = \frac{k^3}{2\pi^2} \mathcal{P}_R(k)$. The end of the slow roll is characterised by the condition that one of the slow roll parameters that are small during inflation becomes $\mathcal{O}(1)$. We obtain the inflaton value at the end of the slow roll φ_* by solving $\max(\epsilon, |\eta|) = 1$, and then find the inflaton value φ_k at the horizon exit of the comoving CMB scale k by solving $N_e = \int_{\varphi_*}^{\varphi_k} d\varphi V_E(d\hat{\varphi}/d\varphi)/(dV_E/d\hat{\varphi})$ for an e-folding number N_e . In this way the value of κ is fixed once N_e , v and ξ are given.

The inflaton potential (18) includes various cases in its limits [24]. When $\xi \rightarrow 0$, $v \rightarrow 0$ it approaches to the minimally coupled $\lambda\phi^4$ model, while the limit $\xi \rightarrow 0$, $v \rightarrow \infty$ gives the prediction obtained in the minimally coupled $m^2\phi^2$ model. It is also known that the $4|\xi|v^2 \rightarrow 1$ limit yields the same inflationary prediction as the nonminimally coupled Higgs inflation with large positive ξ . As an indication of how close to this limit our model is, we introduce a parameter

$$\delta \equiv 1 - 4|\xi|v^2. \quad (19)$$

This is actually Φ of (5) appearing in the Kähler potential, evaluated at the supersymmetric vacuum at $\varphi = 2v$ and $S = 0$. In the limit $\xi \rightarrow 0$ the potential becomes the double-well type; the prediction of this inflationary model is compatible with the combined Planck + WP + BAO results [4] when $2v \gtrsim 13$ in our parametrisation at 95% confidence level (CL). As $|\xi|$ is increased the singularity at $\varphi = 1/\sqrt{|\xi|}$ approaches the potential minimum at $\varphi = 2v$. Fig. 1 shows the shape of the potential when $v = 10$ and $\xi = 0, -0.0010, -0.0015$ and -0.0020 . The scalar spectral index $1 + d \ln A_s(k)/d \ln k \simeq 1 - 6\epsilon_V + 2\eta_V$, the running of the spectral index $dn_s/d \ln k \simeq -24\epsilon_V^2 + 16\epsilon_V \eta_V - 2\xi_V^2$, and the tensor-to-scalar ratio $r \equiv \mathcal{P}_{\text{gw}}/\mathcal{P}_R \simeq 16\epsilon_V$ are evaluated by computing the slow roll parameters at the horizon exit of the comoving CMB scale. We list these results for $v = 10$ and the e-folding number $N_e = 60$ in Table 2, along with the values of κ , φ_* and φ_k found by the procedure explained above. The nonminimal coupling is varied as

⁴ There are FHI models with a noncanonical Kähler potential, including [45–50].

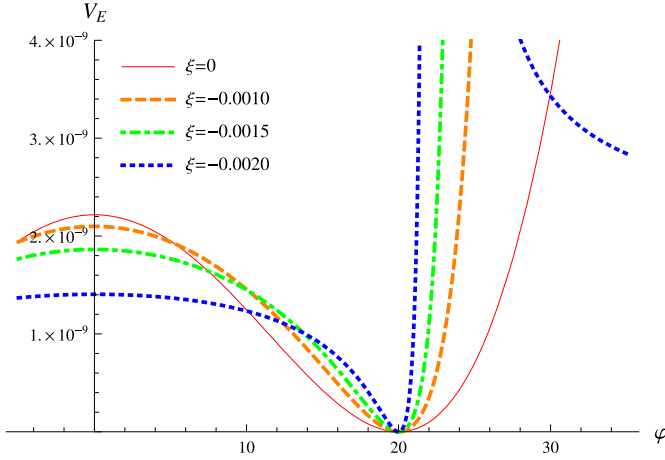


Fig. 1. The scalar potential in the Einstein frame for $\nu = 10$ and $\xi = 0, -0.0010, -0.0015, -0.0020$. The value of κ is determined by the amplitude of the primordial density perturbation (see text). The e-folding number is chosen to be $N_e = 60$.

$\xi = 0, -0.0010, -0.0015, -0.0020$. The tendency of these CMB parameters may be understood from the behaviour of the potential in Fig. 1. As $|\xi|$ is increased, the minimum of the potential becomes a steep valley, while the small ϕ region becomes a plateau; consequently, the spectrum of the inflationary model approaches to that of the nonminimally coupled Higgs inflation model.

4. Comparison with Planck

In this section we show the prediction of our model for the scalar spectral index n_s , the running of the scalar spectral index $dn_s/d \ln k$, and the tensor-to-scalar ratio r for varying ν . The non-minimal coupling is varied as $-1/4\nu^2 < \xi \leq 0$. The procedure and the normalisation are as described in the previous section.

Table 2

The parameter κ , the inflaton field value at the end of the slow roll ϕ_* and at the horizon exit of the comoving CMB scale ϕ_k , the spectral index n_s and its running $dn_s/d \ln k$, the tensor-to-scalar ratio r and the parameter δ defined in (19) computed in the example of $\nu = 10$ and $\xi = 0, -0.0010, -0.0015, -0.0020$. We have chosen the e-folding number $N_e = 60$.

ξ	κ^2	ϕ_*	ϕ_k	n_s	$dn_s/d \ln k$	r	δ
0	2.22×10^{-13}	18.6	7.05	0.964	-4.54×10^{-4}	0.0519	1
-0.0010	2.10×10^{-13}	19.0	9.44	0.968	-4.65×10^{-4}	0.0466	0.6
-0.0015	1.86×10^{-13}	19.2	11.1	0.970	-4.63×10^{-4}	0.0408	0.4
-0.0020	1.40×10^{-13}	19.4	13.5	0.972	-4.56×10^{-4}	0.0313	0.2

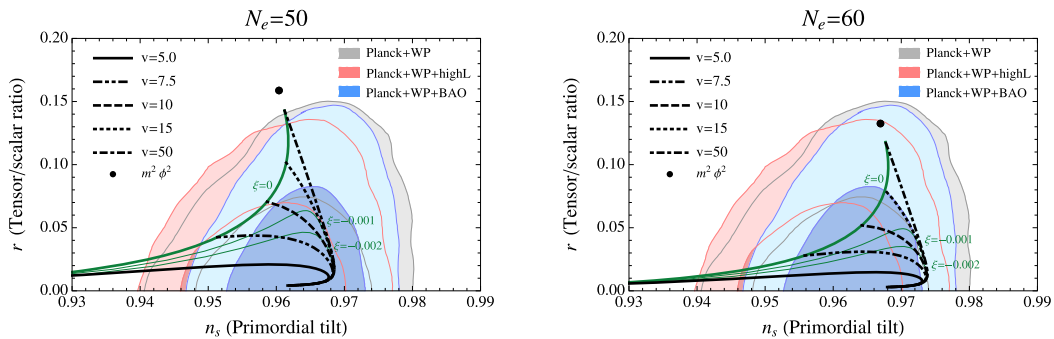


Fig. 2. The spectral index n_s and the tensor-to-scalar ratio r of our model, when the symmetry breaking parameter ν and the value of the nonminimal coupling ξ are varied as $\nu = 5.0, 7.5, 10, 15, 50$ and $-1/4\nu^2 < \xi \leq 0$. The left (right) panel shows the results for the e-folding number $N_e = 50$ ($N_e = 60$). The 68% and 95% CL contours from the Planck satellite observation [3] (Planck + WP: grey, Planck + WP + highL: red, Planck + WP + BAO: blue, from the background to the foreground) are shown for comparison. The black dot is the prediction of the minimally coupled $m^2\phi^2$ chaotic inflation model, i.e. $n_s = 1 - 2/(N_e + \frac{1}{2})$ and $r = 4(1 - n_s)$, to which our model approaches in the limit $\xi = 0, \nu \rightarrow \infty$. The thick green curve is $\xi = 0$ and the thin green curves are $\xi = -0.001$ and $\xi = -0.002$ from above. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

Fig. 2 shows the plots of n_s-r , the left panel showing the results for $N_e = 50$ and the right panel for $N_e = 60$. The 68% and 95% contours from the Planck satellite observation [3] (Planck + WP: grey, Planck + WP + highL: red, Planck + WP + BAO: blue) are superimposed on the background for comparison. In the minimally coupled case ($\xi = 0$) small ν is strongly disfavoured ($\nu \gtrsim 6.5M_{\text{Pl}}$ for Planck + WP + BAO 95% CL [4]). With small negative ξ , in contrast, we see that smaller ν (see $\nu = 5$ for example) is not only compatible but in excellent fit with the current CMB data. This feature is favourable for the model as the large super-Planckian excursion of the inflaton is often considered problematic.

The running of the scalar spectral index $dn_s/d \ln k$ is shown against the scalar spectral index n_s in Fig. 3 (the left panel: $N_e = 50$, the right panel: $N_e = 60$). The 68% and 95% CL contours of the Planck + WP + BAO [4] are also shown for comparison. In the figure the 68% and 95% CL contours of $\Lambda\text{CDM} + dn_s/d \ln k$ are shown by dark and light blue, and the 95% CL contour of $\Lambda\text{CDM} + dn_s/d \ln k + r$ is shown by the light red curve. The 68% CL contour of $\Lambda\text{CDM} + dn_s/d \ln k + r$ is outside the figure. While the running of the scalar spectral index is potentially an important observable beyond n_s and r , the data at present is not significant enough to restrict the model parameters; the contours run nearly vertical in the figure, indicating that the constraints are mainly due to n_s .

Going back to Fig. 2, we see that the prediction of our model for n_s and r makes stark contrast to the nonminimally coupled $\lambda\phi^4$ model (see e.g. [5,22,23]) in which the prediction moves vertically in the n_s-r plane as the nonminimal coupling ξ is varied. In Fig. 2 the parameter space of our model covers almost the whole area inside the 68% CL contour; it would be interesting to see how future observations [56–59], in particular precision measurements of n_s , will constrain these parameters.

5. Reheating and leptogenesis

In this section we discuss viability of the post-inflationary physics. Our scenario is based on the well-motivated particle

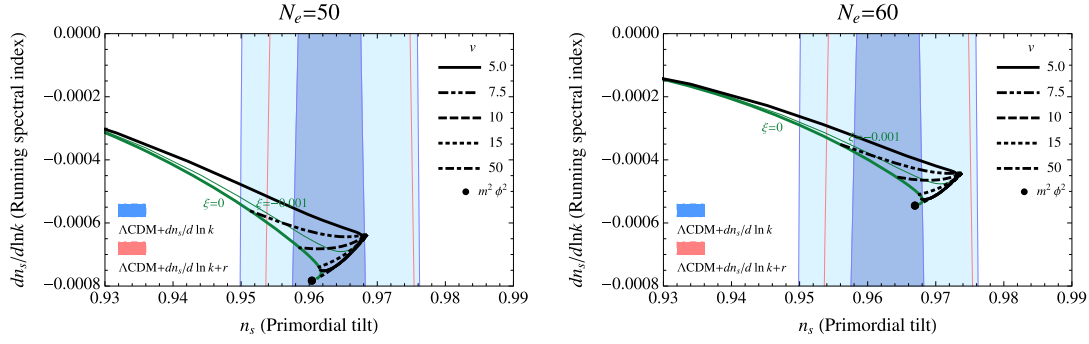


Fig. 3. The prediction of our model for the running of the scalar spectral index $dn_s/d \ln k$ against the scalar spectral index n_s . The left panel shows the results for the e-folding number $N_e = 50$ and the right panel is for $N_e = 60$. The parameters are the same as in Fig. 2. The 68% and 95% CL contours of the Planck + WP + BAO [4] are indicated in the background. The 68% and 95% CL contours of $\Lambda\text{CDM} + dn_s/d \ln k$ are shown by blue and light blue, and the 95% CL contour of $\Lambda\text{CDM} + dn_s/d \ln k + r$ is shown by the thin red curves (the 68% CL contour is outside the figure). The black dot is the prediction of the minimally coupled $m^2 \phi^2$ chaotic inflation model, namely $n_s = 1 - 2/(N_e + \frac{1}{2})$ and $dn_s/d \ln k = -2/(N_e + \frac{1}{2})^2$, to which our model approaches in the $\xi = 0$, $v \rightarrow \infty$ limit. The thick (thin) green curve is the prediction for $\xi = 0$ ($\xi = -0.001$). (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

Inflaton mass (GeV)

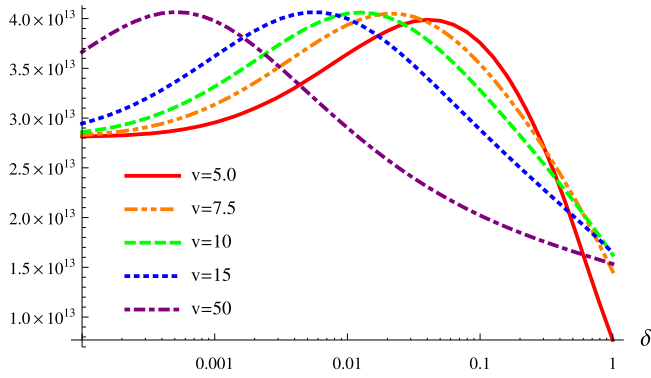


Fig. 4. The inflaton mass $m_{\delta\hat{\phi}}$, plotted against δ (19) for $v = 5.0, 7.5, 10, 15, 50$.

physics model of the gauged $B - L$ extended MSSM with R-symmetry, which has been studied in detail. A peculiar feature of our model is that the breaking scale $v \approx \mathcal{O}(10) \times M_{\text{P}}$ is large, compared to the GUT (see e.g. [38,60]) or the electroweak scale (e.g. [61–64]) $B - L$ breaking scenarios.

Assuming perturbative decay of the inflaton, the upper bound of the reheating temperature is estimated as

$$T_{\text{RH}} \lesssim \left(\frac{90}{g_* \pi^2} \right)^{\frac{1}{4}} \sqrt{M_{\text{P}} \Gamma_{\delta\hat{\phi}}}, \quad (20)$$

where $g_* \approx 200$ is the degrees of freedom at reheating and $\Gamma_{\delta\hat{\phi}}$ is the decay rate of the inflaton in the Einstein frame $\delta\hat{\phi}$ oscillating at the potential minimum. The mass of the inflaton is

$$m_{\delta\hat{\phi}}^2 = \left. \frac{\partial^2 V_{\text{E}}}{\partial \hat{\phi}^2} \right|_{\varphi=2v} = \frac{2\kappa^2 v^2}{1 + 4\xi v^2 (1 + 6\xi)}, \quad (21)$$

which is found to be $m_{\delta\hat{\phi}} = 10^{13} - 10^{14}$ GeV for our model parameters (Fig. 4). We are interested in the decay of the inflaton to the SM particles. The inflaton is a component of Φ_{\pm} and there are two pertinent channels of decay:

- (i) $\delta\hat{\phi} \rightarrow NN, \tilde{N}\tilde{N}$ (the right-handed (s)neutrinos).
- (ii) $\delta\hat{\phi} \rightarrow Z'Z'$ (the $U(1)_{B-L}$ gauge bosons).

Let us consider the decay channel (i) first. The decay in this case is through the coupling $\lambda \equiv \lambda^{ij}$ in the last term of (4). After inflation the field Φ_{-} settles down at $\langle \Phi_{-} \rangle = v$, giving the seesaw scale

$M_N = \lambda(\Phi_{-}) = \lambda v$. Using the neutrino mass $m_{\nu}^2 \approx \Delta m_{32}^2 = 2.43 \times 10^{-3} \text{ eV}^2$ [65] and the Higgs expectation value $\langle H_u \rangle \approx 174 \text{ GeV}$ in the seesaw relation $m_{\nu} M_N = y_D^2 \langle H_u \rangle^2$, we see from $y_D \lesssim \mathcal{O}(1)$ that the seesaw scale is bounded from above: $M_N \lesssim 10^{13} \text{ GeV}$. Since $v \approx \mathcal{O}(10) \times M_{\text{P}}$ in our model, the Majorana Yukawa coupling needs to be small, $\lambda \lesssim 10^{-6}$. The decay rate of the inflaton is

$$\Gamma(\delta\hat{\phi} \rightarrow NN, \tilde{N}\tilde{N}) \approx \frac{\lambda^2}{16\pi} m_{\delta\hat{\phi}}. \quad (22)$$

Using the seesaw relation and $m_{\delta\hat{\phi}} \approx 10^{13} \text{ GeV}$, we find the reheating temperature from (20),

$$T_{\text{RH}} \approx \frac{M_N}{10^{12} \text{ GeV}} \times 10^7 \text{ GeV}. \quad (23)$$

The second channel (ii) also contributes when the Z' mass is smaller than (half of) the inflaton mass $m_{\delta\hat{\phi}}$. From the longitudinal mode dominant decay width

$$\Gamma(\delta\hat{\phi} \rightarrow Z'Z') \approx \frac{1}{32\pi} \frac{m_{\delta\hat{\phi}}^3}{v^2}, \quad (24)$$

the reheating temperature is estimated as

$$T_{\text{RH}} \approx 10^6 \text{ GeV}. \quad (25)$$

The case (i) may be regarded as the dominant channel. The condition that the Big Bang nucleosynthesis is not spoiled by the thermally produced gravitinos yields an upper bound of the reheating temperature $T_{\text{RH}} \lesssim 10^6 - 10^7 \text{ GeV}$ [66,67]. By (23), this gravitino constraint mildly restricts the seesaw scale $M_N \lesssim 10^{12} \text{ GeV}$.

In the out-of-equilibrium decay of the right-handed (s)neutrinos, lepton asymmetry can be generated and is later converted to the baryon asymmetry of the Universe via the sphaleron transitions, the so-called leptogenesis scenario [29,30]. In the sphaleron transition the yield (the ratio of the number density to the entropy density) of the baryons is related to that of the leptons as

$$Y_B \approx -\frac{8}{23} Y_L. \quad (26)$$

In *thermal leptogenesis* scenario in which the right-handed (s)neutrinos are thermally produced, the reheating temperature needs to be higher than the mass scale of the right-handed (s)neutrinos: $T_{\text{RH}} \gtrsim M_N$. The generated baryon asymmetry is estimated as [68–70]

$$Y_B \sim Y_L \sim a \frac{\varepsilon_i}{g_*}, \quad (27)$$

where ε_i is the CP asymmetry parameter associated with the i 'th generation of the right-handed (s)neutrino N_i (\tilde{N}_i), and $a \lesssim 1$ is the efficiency factor which depends on details of the Boltzmann equations. The baryon asymmetry of the Universe is observed to be [2,71]

$$Y_B = (8.55 \pm 0.217) \times 10^{-11} \quad (95\% \text{ CL}), \quad (28)$$

and thus the CP asymmetry parameter needs to be $\varepsilon_i \approx 10^{-7}$. This condition is known to be satisfied when the right-handed (s)neutrino mass is large enough, $M_{N_i} \gtrsim 10^{10}$ GeV [68,69]. In our scenario, however, this requirement cannot be fulfilled as the reheating temperature (23), (25) is not high enough to produce such heavy right-handed (s)neutrinos. Nevertheless, it is known that even if the right-handed (s)neutrino masses are not large, large enough ε_i (and thus Y_B) can be obtained when at least two of the right-handed (s)neutrino masses are nearly degenerate and resonant enhancement takes place [72,73]. Thus, we may think of the following two possible cases of thermal leptogenesis in our scenario: (a) when the decay channel (i) is dominant, the reheating temperature is given by (23) and thus one of the right-handed (s)neutrino masses needs to be large. For successful resonant thermal leptogenesis the remaining two masses need to be very close to each other and less than the reheating temperature, for example, $M_{N_1} \approx M_{N_2} \ll T_{\text{RH}} \ll M_{N_3}$; (b) when the reheating temperature is determined by the decay channel (ii) and is given by (25), all N_i 's can be light: $M_{N_1}, M_{N_2}, M_{N_3} \ll T_{\text{RH}}$. In both (a) and (b), the right-handed (s)neutrinos that produce lepton number through the decay are much lighter than 10^{10} GeV and hence the resonance enhancement needs to take place. For detailed analysis of resonant leptogenesis in the context of the minimal $B-L$ model, see, for example [63,74].

In *nonthermal leptogenesis* scenario, the seesaw scale M_N is larger than the reheating temperature T_{RH} and the right-handed (s)neutrinos are predominantly produced by the decay process (i). The baryon asymmetry generated by the decay of the (s)neutrinos may be estimated as [70]

$$Y_B \sim Y_L \sim \frac{T_{\text{RH}}}{m_{\delta\hat{\varphi}}} \sum_i \text{Br}_i \varepsilon_i, \quad (29)$$

where

$$\text{Br}_i = \frac{\Gamma(\delta\hat{\varphi} \rightarrow N_i N_i, \tilde{N}_i \tilde{N}_i)}{\Gamma_{\delta\hat{\varphi}}} \quad (30)$$

is the branching ratio and $\Gamma_{\delta\hat{\varphi}}$ is the total decay width of the inflaton. Using $T_{\text{RH}} \sim 10^7$ GeV and $m_{\delta\hat{\varphi}} \approx 10^{13}$ GeV in (29) we find $\text{Br}_i \varepsilon_i \approx 10^{-4}$. This large CP asymmetry parameter can be again obtained by resonant leptogenesis. To conclude, for the production of the baryon asymmetry through thermal or nonthermal leptogenesis the CP asymmetry parameter needs to be $\varepsilon_i \sim 10^{-7}$ – 10^{-4} . These values are somewhat larger than the conventional decay scenarios of the right-handed (s)neutrinos, but can be accounted for by the resonance enhancement [72,73]. For this, at least two of the right-handed neutrino masses need to be nearly degenerate. The baryon asymmetry may, alternatively, be generated by some other mechanism such as the Affleck–Dine mechanism [75].

6. Discussions

In this paper we have constructed a novel scenario of cosmological inflation based on the gauged $B-L$ extended MSSM. The model is well-motivated by the neutrino physics: the spontaneous breaking of the $U(1)_{B-L}$ symmetry gives rise to the right-handed neutrino mass term, which in turn give the small nonzero left-handed neutrino masses through the type I seesaw mechanism. It

also includes the mechanism of baryogenesis through leptogenesis. Due to supersymmetry our model is stable against radiative corrections and includes supersymmetric particles that may be considered a good dark matter candidate. Our model has various advantages over the popular FHI models. For example the observationally supported slightly red spectrum of the primordial density perturbation can be naturally accounted for. There is no need to invoke radiative corrections, and thus there is no necessity of fine-tuning associated with it.

A notable feature of our scenario is that it is free from unwanted relic particles. As discussed in Section 5 the reheating temperature is low enough so that the gravitino problem can be avoided. In addition, our model is free from cosmic strings that generally impose stringent constraints on the FHI models (see e.g. [33]). This is due to our assumption that the $U(1)_{B-L}$ symmetry is already broken at the onset of inflation and cosmic strings are inflated away.

We conclude the paper by commenting on issues that are of potential importance but were not discussed above. The original model (3) includes multiple fields and we simplified the model by focusing on the φ direction. This is justified by assuming the quartic term in the Kähler potential that controls the tachyonic instability (see [17,18]). Lifting this assumption certainly complicates the scenario, but leads to rich observable consequences. The resulting model is sensitive to the initial condition of the inflaton field, as in the case of the FHI scenario. The extra degrees of freedom give rise to the isocurvature mode and possibly non-Gaussianity of the density perturbation (see e.g. [76] for a review). If such a signature is to be detected in the future, it will be an indication that the single field approximation that we used is clearly inappropriate. Our second comment concerns the interpretation of the $\delta \rightarrow 0$ limit. The δ defined in (19) is the factor Φ of the Kähler potential (5) at the potential minimum. There is no reason for it to be extremely large or small, and thus the ‘‘Higgs-inflation limit’’ $\delta \rightarrow 0$ is a limit of fine-tuning. In this sense, this $\delta \rightarrow 0$ limit is not much better than the $\xi \gg 1$ Higgs inflation. In Fig. 2, $\delta = 1$ ($\xi = 0$) is indicated by the thick green curve and smaller δ (larger $|\xi|$) corresponds to smaller r . See also Table 2. The CMB polarisation experiments [56–59] are expected to uncover the physics of primordial gravitational waves, with accuracy corresponding to $r \sim 0.01$. These experiments will tell us whether one actually needs to consider the fine-tuned small δ limit. Finally we comment on possible extensions of the model. The essential elements of our model are the superpotential (3) and the Kähler potential (5), and hence, it is easy to construct a similar inflaton potential if a supersymmetric model is equipped with the same structure. It is, however, not straightforward to construct a consistent scenario of inflation since keeping the R-symmetry in a realistic GUT is known to be extremely difficult [77].

7. Note added

After completion of this paper detection of the CMB B-mode polarisation was announced by the BICEP2 experiment [78], with the tensor/scalar ratio $r = 0.20^{+0.07}_{-0.05}$. This means in our model that small values of $|\xi|$, and hence $\delta \sim 1$ *without* fine-tuning (see (19)), are favoured.

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References

- [1] C. Bennett, et al., *Astrophys. J. Suppl. Ser.* 208 (2013) 20, arXiv:1212.5225 [astro-ph.CO].
- [2] G. Hinshaw, et al., *Astrophys. J. Suppl. Ser.* 208 (2013) 19, arXiv:1212.5226 [astro-ph.CO].
- [3] P. Ade, et al., Planck Collaboration, arXiv:1303.5062 [astro-ph.CO], 2013.
- [4] P. Ade, et al., Planck Collaboration, arXiv:1303.5082 [astro-ph.CO], 2013.
- [5] N. Okada, M.U. Rehman, Q. Shafi, *Phys. Rev. D* 82 (2010) 043502, arXiv:1005.5161 [hep-ph].
- [6] J. Cervantes-Cota, H. Dehnen, *Nucl. Phys. B* 442 (1995) 391, arXiv:astro-ph/9505069.
- [7] F. Bezrukov, M. Shaposhnikov, *Phys. Lett. B* 659 (2008) 703, arXiv:0710.3755 [hep-th].
- [8] A. De Simone, M.P. Hertzberg, F. Wilczek, *Phys. Lett. B* 678 (2009) 1, arXiv:0812.4946 [hep-ph].
- [9] F. Bezrukov, M. Shaposhnikov, *J. High Energy Phys.* 0907 (2009) 089, arXiv:0904.1537 [hep-ph].
- [10] A. Barvinsky, A.Y. Kamenshchik, C. Kiefer, A. Starobinsky, C. Steinwachs, *J. Cosmol. Astropart. Phys.* 0912 (2009) 003, arXiv:0904.1698 [hep-ph].
- [11] C. Burgess, H.M. Lee, M. Trott, *J. High Energy Phys.* 0909 (2009) 103, arXiv:0902.4465 [hep-ph].
- [12] J. Barbon, J. Espinosa, *Phys. Rev. D* 79 (2009) 081302, arXiv:0903.0355 [hep-ph].
- [13] C. Burgess, H.M. Lee, M. Trott, *J. High Energy Phys.* 1007 (2010) 007, arXiv:1002.2730 [hep-ph].
- [14] F. Bezrukov, M.Y. Kalmykov, B.A. Knieh, M. Shaposhnikov, *J. High Energy Phys.* 1210 (2012) 140, arXiv:1205.2893 [hep-ph].
- [15] A. Salvio, *Phys. Lett. B* 727 (2013) 234, arXiv:1308.2244 [hep-ph].
- [16] M.B. Einhorn, D.T. Jones, *J. High Energy Phys.* 1003 (2010) 026, arXiv:0912.2718 [hep-ph].
- [17] S. Ferrara, R. Kallosh, A. Linde, A. Marrani, A. Van Proeyen, *Phys. Rev. D* 82 (2010) 045003, arXiv:1004.0712 [hep-th].
- [18] S. Ferrara, R. Kallosh, A. Linde, A. Marrani, A. Van Proeyen, *Phys. Rev. D* 83 (2011) 025008, arXiv:1008.2942 [hep-th].
- [19] M. Arai, S. Kawai, N. Okada, *Phys. Rev. D* 84 (2011) 123515, arXiv:1107.4767 [hep-ph].
- [20] C. Pallis, N. Toumbas, *J. Cosmol. Astropart. Phys.* 1112 (2011) 002, arXiv:1108.1771 [hep-ph].
- [21] M.B. Einhorn, D.T. Jones, *J. Cosmol. Astropart. Phys.* 1211 (2012) 049, arXiv:1207.1710 [hep-ph].
- [22] M. Arai, S. Kawai, N. Okada, *Phys. Rev. D* 86 (2012) 063507, arXiv:1112.2391 [hep-ph].
- [23] M. Arai, S. Kawai, N. Okada, *Phys. Rev. D* 87 (2013) 065009, arXiv:1212.6828 [hep-ph].
- [24] A. Linde, M. Noorbala, A. Westphal, *J. Cosmol. Astropart. Phys.* 1103 (2011) 013, arXiv:1101.2652 [hep-th].
- [25] R. Kallosh, A. Linde, *J. Cosmol. Astropart. Phys.* 1307 (2013) 002, arXiv:1306.5220 [hep-th].
- [26] M.S. Carena, A. Daleo, B.A. Dobrescu, T.M. Tait, *Phys. Rev. D* 70 (2004) 093009, arXiv:hep-ph/0408098.
- [27] G. Cacciapaglia, C. Csaki, G. Marandella, A. Strumia, *Phys. Rev. D* 74 (2006) 033011, arXiv:hep-ph/0604111.
- [28] P. Minkowski, *Phys. Lett. B* 67 (1977) 421; T. Yanagida, in: O. Sawada, A. Sugamoto (Eds.), *Proc. of the Workshop on the Baryon Number of the Universe and Unified Theories*, Tsukuba, Japan, 13–14 February 1979, 1979, p. 95, KEK report KEK-79-18; M. Gell-Mann, P. Ramond, R. Slansky, in: P. van Nieuwenhuizen, D.Z. Freedman (Eds.), *Supergravity*, North Holland Publ. Co, 1979, pp. 315–321, 1979, print-80-0576 (CERN); R.N. Mohapatra, G. Senjanovic, *Phys. Rev. Lett.* 44 (1980) 912.
- [29] M. Fukugita, T. Yanagida, *Phys. Lett. B* 174 (1986) 45.
- [30] G. Lazarides, Q. Shafi, *Phys. Lett. B* 258 (1991) 305.
- [31] E.J. Copeland, A.R. Liddle, D.H. Lyth, E.D. Stewart, D. Wands, *Phys. Rev. D* 49 (1994) 6410, arXiv:astro-ph/9401011.
- [32] G. Dvali, Q. Shafi, R.K. Schaefer, *Phys. Rev. Lett.* 73 (1994) 1886, arXiv:hep-ph/9406319.
- [33] G. Lazarides, C. Panagiotakopoulos, *Phys. Rev. D* 52 (1995) 559, arXiv:hep-ph/9506325.
- [34] G. Lazarides, R.K. Schaefer, Q. Shafi, *Phys. Rev. D* 56 (1997) 1324, arXiv:hep-ph/9608256.
- [35] R. Jeannerot, arXiv:hep-ph/9802332, 1997.
- [36] R. Jeannerot, S. Khalil, G. Lazarides, Q. Shafi, *J. High Energy Phys.* 0010 (2000) 012, arXiv:hep-ph/0002151.
- [37] R. Jeannerot, S. Khalil, G. Lazarides, *Phys. Lett. B* 506 (2001) 344, arXiv:hep-ph/0103229.
- [38] V.N. Senoguz, Q. Shafi, arXiv:hep-ph/0512170, 2005.
- [39] M.U. Rehman, Q. Shafi, J.R. Wickman, *Phys. Lett. B* 683 (2010) 191, arXiv:0908.3896 [hep-ph].
- [40] W. Buchmuller, V. Domcke, K. Kamada, K. Schmitz, *J. Cosmol. Astropart. Phys.* 1310 (2013) 003, arXiv:1305.3392 [hep-ph].
- [41] N. Okada, M.U. Rehman, Q. Shafi, *Phys. Lett. B* 701 (2011) 520, arXiv:1102.4747 [hep-ph].
- [42] N. Okada, Q. Shafi, arXiv:1311.0921 [hep-ph], 2013.
- [43] G. Lazarides, Q. Shafi, *Phys. Rev. D* 58 (1998) 071702, arXiv:hep-ph/9803397.
- [44] M. Kaku, P.K. Townsend, P. van Nieuwenhuizen, *Phys. Rev. D* 17 (1978) 3179; W. Siegel, S.J. Gates Jr., *Nucl. Phys. B* 147 (1979) 77; E. Cremmer, S. Ferrara, L. Girardello, A. Van Proeyen, *Nucl. Phys. B* 212 (1983) 413; S. Ferrara, L. Girardello, T. Kugo, A. Van Proeyen, *Nucl. Phys. B* 223 (1983) 191; T. Kugo, S. Uehara, *Nucl. Phys. B* 222 (1983) 125; T. Kugo, S. Uehara, *Nucl. Phys. B* 226 (1983) 49; T. Kugo, S. Uehara, *Prog. Theor. Phys.* 73 (1985) 235.
- [45] M. Bastero-Gil, S. King, Q. Shafi, *Phys. Lett. B* 651 (2007) 345, arXiv:hep-ph/0604198.
- [46] B. Garbrecht, C. Pallis, A. Pilaftsis, *J. High Energy Phys.* 0612 (2006) 038, arXiv:hep-ph/0605264.
- [47] M. ur Rehman, V.N. Senoguz, Q. Shafi, *Phys. Rev. D* 75 (2007) 043522, arXiv:hep-ph/0612023.
- [48] C. Pallis, *J. Cosmol. Astropart. Phys.* 0904 (2009) 024, arXiv:0902.0334 [hep-ph].
- [49] M.U. Rehman, Q. Shafi, J.R. Wickman, *Phys. Rev. D* 83 (2011) 067304, arXiv:1012.0309 [astro-ph.CO].
- [50] R. Armillius, C. Pallis, in: A. Travena, B. Soren (Eds.), *Recent Advances in Cosmology*, Nova Science Publishers, 2013, p. 159, arXiv:1211.4011 [hep-ph].
- [51] C. Pallis, Q. Shafi, *Phys. Lett. B* 725 (2013) 327, arXiv:1304.5202 [hep-ph].
- [52] S. Orani, A. Rajantie, *Phys. Rev. D* 88 (2013) 043508, arXiv:1304.8041 [astro-ph.CO].
- [53] L. Boubekeur, D. Lyth, *J. Cosmol. Astropart. Phys.* 0507 (2005) 010, arXiv:hep-ph/0502047.
- [54] A.D. Linde, *Pis'ma Zh. Eksp. Teor. Fiz.* 30 (1979) 479.
- [55] A.A. Starobinsky, *Sov. Astron. Lett.* 7 (1981) 36.
- [56] M. Hazumi, *AIP Conf. Proc.* 1040 (2008) 78.
- [57] J. Bock, et al., EPIC Collaboration, arXiv:0906.1188 [astro-ph.CO], 2009.
- [58] F. Bouchet, et al., CoRE Collaboration, arXiv:1102.2181 [astro-ph.CO], 2011.
- [59] Z. Kermish, P. Ade, A. Anthony, K. Arnold, K. Arnold, et al., arXiv:1210.7768 [astro-ph.IM], 2012.
- [60] W. Buchmuller, V. Domcke, K. Schmitz, *Nucl. Phys. B* 862 (2012) 587, arXiv:1202.6679 [hep-ph].
- [61] S. Iso, N. Okada, Y. Orikasa, *Phys. Lett. B* 676 (2009) 81, arXiv:0902.4050 [hep-ph].
- [62] S. Iso, N. Okada, Y. Orikasa, *Phys. Rev. D* 80 (2009) 115007, arXiv:0909.0128 [hep-ph].
- [63] S. Iso, N. Okada, Y. Orikasa, *Phys. Rev. D* 83 (2011) 093011, arXiv:1011.4769 [hep-ph].
- [64] Z.M. Burell, N. Okada, *Phys. Rev. D* 85 (2012) 055011, arXiv:1111.1789 [hep-ph].
- [65] G. Fogli, E. Lisi, A. Marrone, D. Montanino, A. Palazzo, et al., *Phys. Rev. D* 86 (2012) 013012, arXiv:1205.5254 [hep-ph].
- [66] M. Kawasaki, K. Kohri, T. Moroi, *Phys. Lett. B* 625 (2005) 7, arXiv:astro-ph/0402490.
- [67] M. Kawasaki, K. Kohri, T. Moroi, *Phys. Rev. D* 71 (2005) 083502, arXiv:astro-ph/0408426.
- [68] W. Buchmuller, P. Di Bari, M. Plumacher, *Nucl. Phys. B* 643 (2002) 367, arXiv:hep-ph/0205349.
- [69] W. Buchmuller, P. Di Bari, M. Plumacher, *Ann. Phys.* 315 (2005) 305, arXiv:hep-ph/0401240.
- [70] W. Buchmuller, R. Peccei, T. Yanagida, *Annu. Rev. Nucl. Part. Sci.* 55 (2005) 311, arXiv:hep-ph/0502169.
- [71] P. Ade, et al., Planck Collaboration, arXiv:1303.5076 [astro-ph.CO], 2013.
- [72] M. Flanz, E.A. Paschos, U. Sarkar, J. Weiss, *Phys. Lett. B* 389 (1996) 693, arXiv:hep-ph/9607310.
- [73] A. Pilaftsis, *Phys. Rev. D* 56 (1997) 5431, arXiv:hep-ph/9707235.
- [74] N. Okada, Y. Orikasa, T. Yamada, *Phys. Rev. D* 86 (2012) 076003, arXiv:1207.1510 [hep-ph].
- [75] I. Affleck, M. Dine, *Nucl. Phys. B* 249 (1985) 361.
- [76] A. Mazumdar, J. Rocher, *Phys. Rep.* 497 (2011) 85, arXiv:1001.0993 [hep-ph].
- [77] M. Fallbacher, M. Ratz, P.K. Vaudrevange, *Phys. Lett. B* 705 (2011) 503, arXiv:1109.4797 [hep-ph].
- [78] P. Ade, et al., BICEP2 Collaboration, arXiv:1403.3985 [astro-ph.CO], 2014.