# $U(1)_{B-L}$ extra-natural inflation with Standard Model on a brane

Kazuyuki Furuuchi [a,*], Jackson M.S. Wu [b]

[a] Manipal Center for Natural Sciences, Manipal University, Madhav Nagar, Manipal, Karnataka 576104, India
[b] Physics Division, National Center for Theoretical Sciences, No. 101, Section 2, Kuang Fu Road, Hsinchu 30013, Taiwan, ROC

## ARTICLE INFO

## ABSTRACT

The interrelation between inflationary cosmology and new physics beyond the Standard Model (SM) is studied in a $U(1)_{B-L}$ extension of the SM embedded in a $(4+1)$-dimensional spacetime. In the scenario we study, the inflaton arises from the Wilson loop of the $U(1)_{B-L}$ gauge group winding an extra-dimensional cycle. Particular attention is paid to the coupling between the inflaton and SM particles that are confined on a brane localized in the extra dimension. We find that the inflaton decay channels are rather restricted in this scenario and the resulting reheating temperature is relatively low.

## 1. Introduction

The precision of the Cosmic Microwave Background (CMB) anisotropy observations has started to rule out some of the inflation models [1]. However, CMB data alone still accommodates a large class of them. In order to narrow down further likely candidates, it is useful to study possible relevance of the inflaton to physics in other eras. In particular, at the time of reheating, the inflaton decays to Standard Model (SM) particles so that the standard hot big bang can proceed, the nature of the interaction between the inflaton and the SM is thus crucial.

Large field inflation models had attracted attention because of the possible detection of tensor modes in CMB polarization in the near future [2]. It is theoretically challenging to construct natural large field inflation models, since effective field theory approach usually breaks down in these models. Extra-natural inflation [3,4], which is based on a gauge theory in higher-dimensional spacetime, is one way to circumvent this difficulty by using non-local operator (Wilson loop) in the extra dimension.

It is an interesting question what should be the gauge group for extra-natural inflation. As we review in the next section, it turns out that to explain the CMB data the gauge coupling for extra-natural inflation must be very small [3,4]. This makes it difficult to identify the SM gauge groups as that for extra-natural inflation, as their couplings at the electro-weak scale are orders of magnitudes larger than that required for extra-natural inflation. Therefore we shall look for other gauge groups in models beyond the SM (BSM).

Gauged $U(1)_{B-L}$ extension of the SM [5–8] is ubiquitous in scenarios of BSM physics. A nice feature of it is that the existence of

right-handed neutrinos is made natural by the necessity of gauge anomaly cancellation. It also makes $R$-parity exact in supersymmetric versions of the SM, and it appears as an intermediate stage in the symmetry breaking pattern of grand-unified models down to the SM, as well as in higher-dimensional embeddings of the SM in string theory constructions. Apart from the formal theoretical considerations, phenomenologically, having a new gauge boson and scalars neutral under the SM gauge group can give rise to novel effects observable in future collider experiments.

In this Letter, we study extra-natural inflation with $U(1)_{B-L}$ as the gauge group. In the scenario we study, the bulk spacetime is $(4+1)$-dimensional with the extra dimension compactified on a circle, SM is confined on a $(3+1)$-dimensional brane localized in the extra dimension, and the inflation arises from the Wilson loop of the $U(1)_{B-L}$ gauge field living in the full $(4+1)$-dimensional bulk. In the following, we explore the interrelation between inflationary cosmology and particle physics in this setting.[1]

The rest of the Letter is organized as follows. The relevant ingredients of extra-natural inflation is reviewed in Section 2. The details of our $U(1)_{B-L}$ extension of the SM is discussed in Section 3. The decay of the inflaton to SM particles is studied in Section 4. We end with a summary and discussions in Section 5.

## 2. $U(1)_{B-L}$ extra-natural inflation

Extra-natural inflation [3,4] is a version of natural inflation [11] whose typical potential takes the form

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* Corresponding author.

[1] For other approaches to connect inflation and new physics beyond SM via extra-natural inflaton, see [9,10].

$$V(\phi) = \frac{V_0}{2}\left[1 - \cos\left(\frac{\phi}{f}\right)\right], \tag{2.1}$$

where $\phi$ is the inflaton which, in extra-natural inflation, is the zero-mode of the fifth component of some bulk gauge field. In the scenario we study here, it is that of the $U(1)_{B-L}$ gauge group. From (2.1) the slow-roll parameters are given by

$$\epsilon_V(\phi) \equiv \frac{M_P^2}{2}\left(\frac{V'}{V}\right)^2 = \frac{M_P^2}{2f^2}\left(\frac{\sin(\frac{\phi}{f})}{1 - \cos(\frac{\phi}{f})}\right)^2, \tag{2.2}$$

$$\eta_V(\phi) \equiv M_P^2 \frac{V''}{V} = \frac{M_P^2}{f^2}\frac{\cos(\frac{\phi}{f})}{1 - \cos(\frac{\phi}{f})}. \tag{2.3}$$

Here $'$ denotes derivative with respect to $\phi$. The slow-roll conditions amount to

$$\epsilon_V, |\eta_V| \ll 1. \tag{2.4}$$

In extra-natural inflation, $f$ and $V_0$ are estimated as [3]

$$f = \frac{1}{g_4(2\pi L_5)}, \tag{2.5}$$

and

$$V_0 = \frac{c_0}{\pi^2}\frac{1}{(2\pi L_5)^4}. \tag{2.6}$$

Here, $g_4$ is the (effective) four-dimensional gauge coupling, and $L_5$ is the radius of the compactified fifth dimension. The constant $c_0$ is determined by the matter content in the bulk, with the relevant ones being fields charged under the gauge symmetry of interest and whose masses are below or of the order of $1/L_5$ [12]; each of these field makes an $\mathcal{O}(1)$ contribution to $c_0$.[2]

In order for quantum gravity corrections to be small, we need

$$(L_5 M_5)^3 \gg 1, \tag{2.7}$$

where $M_5$ is the five-dimensional (reduced) Planck scale, which is related to the four-dimensional reduced Planck scale $M_P \simeq 2.4 \times 10^{18}$ GeV by

$$M_P^2 = M_5^3 (2\pi L_5). \tag{2.8}$$

Thus from (2.5)

$$M_5 = (g_4 f M_P^2)^{1/3}. \tag{2.9}$$

Since $f$ is directly related to the CMB observations, and $g_4$ is a basic parameter in the $U(1)_{B-L}$ extension of the SM, we shall take $f$ and $g_4$ as the independent parameters, and regard $L_5$ and $M_5$ as functions of them. It is convenient to introduce a dimensionless parameter

$$\ell_5 \equiv L_5 M_5 = \frac{1}{2\pi}\left(\frac{M_P}{g_4 f}\right)^{2/3}, \tag{2.10}$$

which measures the strength of quantum gravity corrections; (2.7) then amounts to $\ell_5^3 \gg 1$. Although $\ell_5$ is not an independent parameter, it is sometimes convenient to use $\ell_5$ instead of $g_4$. In terms of $\ell_5$ and $f$, $g_4$ is expressed as

$$g_4 = \frac{M_P}{f}(2\pi \ell_5)^{-3/2}. \tag{2.11}$$

The number of e-folds as a function of $\phi$ is given by

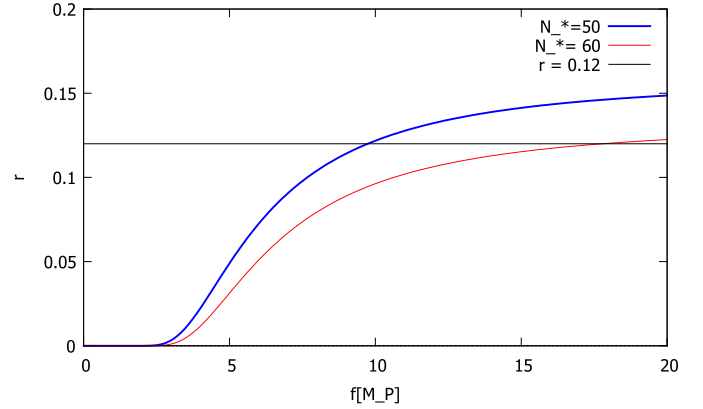

Fig. 1. The scalar-to-tensor ratio, $r$, as a function of $f$ for different values of $N_*$.

$$N(\phi) = \int_t^{t_e} dt\, H(\phi(t)) = \int_\phi^{\phi_e} d\phi\, \frac{H}{\dot\phi}$$

$$\simeq -\int_\phi^{\phi_e} d\phi\, \frac{1}{M_P^2}\frac{V}{V'} = \int_{\phi_e}^\phi d\phi\, \frac{f}{M_P^2}\frac{1 - \cos\frac{\phi}{f}}{\sin\frac{\phi}{f}}$$

$$= -\left(\frac{f}{M_P}\right)^2 \log\left[\frac{1}{2}\left(1 + \cos\frac{\phi}{f}\right)\right]_{\phi_e}^\phi. \tag{2.12}$$

Here, $\phi_e$ is the value of the inflaton field at the end of inflation defined by $\epsilon_V(\phi_e) = 1$, where the slow-roll condition (2.4) breaks down.[3] This gives

$$\frac{\phi_e}{f} = \cos^{-1}\left(\frac{1 - \frac{M_P^2}{2f^2}}{1 + \frac{M_P^2}{2f^2}}\right), \tag{2.13}$$

and plugging (2.13) into (2.12) we obtain

$$\frac{\phi}{f} = \cos^{-1}\left(\frac{2e^{-\frac{M_P^2}{f^2}N}}{1 + \frac{M_P^2}{2f^2}} - 1\right). \tag{2.14}$$

In slow-roll inflation, the tensor-to-scalar ratio, $r$, and the spectral index, $n_s$, are given by

$$r \simeq 16\epsilon_V, \qquad n_s \simeq 1 - 6\epsilon_V + 2\eta_V. \tag{2.15}$$

The scalar-to-tensor ratio and the spectral index estimated from various combinations of the Planck data and other observations give at 95% CL: $r \lesssim 0.12$ and $0.94 \lesssim n_s \lesssim 0.98$ at the pivot scale $k_* = 0.002$ Mpc$^{-1}$ [1]. Below, except for $r$ and $n_s$ whose value we take always at the pivot scale, we shall use the subscript $*$ to indicate that the value is taken at the pivot scale.

As can be seen from (2.2), (2.3) and (2.14), $r$ and $n_s$ only depend on $f$ and $N_*$ in extra-natural inflation, and so constraints on $r$ and $n_s$ constrain $f$ for a given $N_*$. We plot the dependence of $r$ and $n_s$ on $f$ at fixed values of $N_*$ in Figs. 1 and 2, respectively. We see that for $N_* = 50$, we have $f \lesssim 10 M_P$ from $r \lesssim 0.12$ and $f \gtrsim 5 M_P$ from $n_s \gtrsim 0.94$. We will see later when considering the inflaton decay that $N_* \simeq 50$ is natural for the scenario we study here.

The power spectrum of the slow-roll inflation is given by

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[2] More precisely, we assume that contributions from charge one fields dominate, which gives rise to the periodicity $\phi \sim \phi + 2\pi f$.

[3] Note that $\epsilon_V \geqslant |\eta_V|$ for $f > M_P$, which is the case in the following.
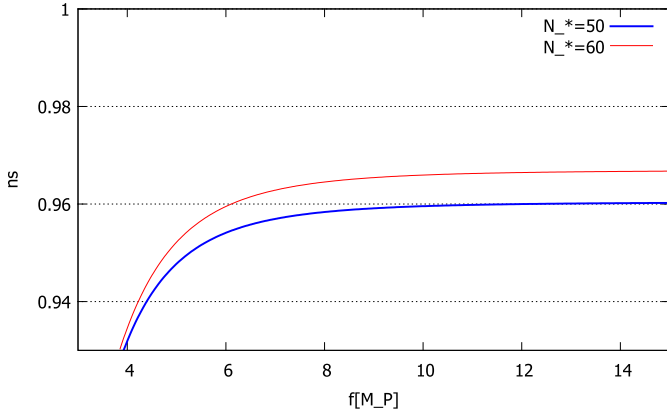
**Fig. 2.** The spectral index, $n_s$, as a function of $f$ for different values of $N_*$.
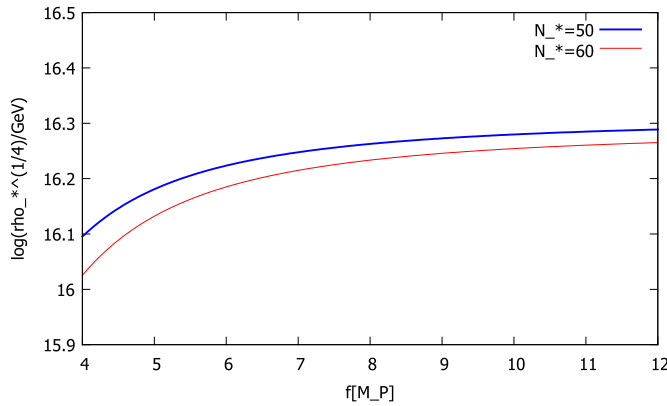


**Fig. 3.** The energy density, $\rho_*$, as a function of $f$ for different values of $N_*$.
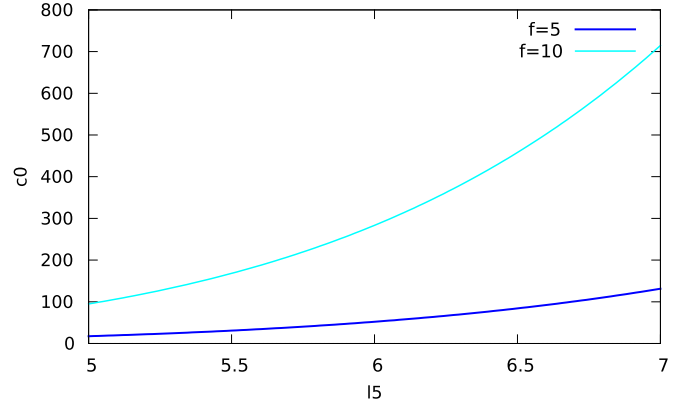


**Fig. 4.** The behavior of $c_0$ as a function of $\ell_5$ for $N_* = 50$ at different values of $f$.
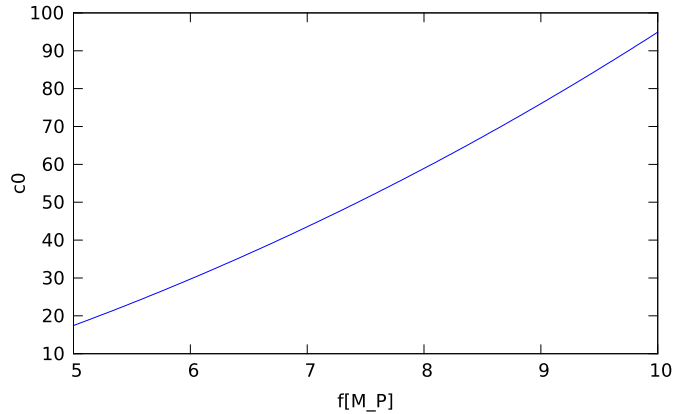


**Fig. 5.** The behavior of $c_0$ as a function of $f$ for $N_* = 50$ and $\ell_5 = 5$.

$$P_\zeta \simeq \frac{H^2}{8\pi^2 M_P^2 \epsilon_V}. \tag{2.16}$$

This should be compared with the observed value $P_\zeta(k_*) = 2.2 \times 10^{-9}$ [1]. It determines the Hubble scale, $H_*$, when the pivot scale exited the horizon, and thus the energy density at that time, $\rho_* \simeq 3M_P^2 H_*^2$, as a function of $f$ and $N_*$. Its dependence on $f$ and $N_*$ is mild, and we obtain $\rho_* \simeq 10^{16}$ GeV, see Fig. 3. On the other hand, from the Friedman equation for spatially flat Universe in the slow-roll approximation,

$$3H^2 M_P^2 = \rho \simeq V(\phi), \tag{2.17}$$

we obtain

$$\begin{aligned}
P_\zeta(k_*) \\
\simeq \frac{V(\phi_*)}{24\pi^2 M_P^4 \epsilon_{V*}} \\
= \frac{c_0}{2\pi^2} \frac{1}{(2\pi L_5)^4} \left[ 1 - \cos\left(\frac{\phi_*}{f}\right) \right] \frac{1}{24\pi^2 M_P^4 \epsilon_{V*}} \\
= \frac{c_0}{2\pi^2} \frac{1}{(2\pi \ell_5)^6} \left[ 1 - \cos\left(\frac{\phi(f, N_*)}{f}\right) \right] \frac{1}{24\pi^2 \epsilon_V(\phi(f, N_*))}.
\end{aligned} \tag{2.18}$$

In the last line we have made it explicit that $\phi$ and $\epsilon_V$ are functions of $f$ and $N$. Thus given $\ell_5$, $f$ and $N_*$, $c_0$ is determined from the observed value $P_\zeta(k_*) = 2.2 \times 10^{-9}$ by (2.18). The behavior of $c_0$ as a function of $\ell_5$ is plotted in Fig. 4. We observe that $c_0$ grows as $\ell_5^6$. Also, $c_0$ grows rapidly with $f$, as seen in Fig. 5. Since

each field charged under $U(1)_{B-L}$ with mass $\lesssim L_5^{-1}$ makes an $\mathcal{O}(1)$ contribution to $c_0$, if it is much larger than unity it may not be natural.[4] Therefore we regard smaller values of $\ell_5$ and $f$, viz. $\ell_5 \simeq 5$ and $f \simeq 5$, as preferred in our scenario here. With the independent parameters fixed, we then have $g_4 \simeq 10^{-3}$ and $L_5^{-1} \simeq 9 \times 10^{16}$ GeV from (2.11) and (2.9). The value of the $U(1)_{B-L}$ gauge coupling is an important input to our $U(1)_{B-L}$ extension of the SM, which we discuss next.

## 3. $U(1)_{B-L}$ extension of the Standard Model

There are several possibilities for the $U(1)_{B-L}$ extension of the SM, particularly with regards to the charge assignment of the scalar field that would break the $U(1)_{B-L}$ symmetry. Table 1 lists the particle content and the charge assignments of the particular $U(1)_{B-L}$ extension of the SM we consider here. In our set-up, we envisage all the SM particles and the right-handed neutrinos living on a four-dimensional brane, while the $U(1)_{B-L}$ gauge fields, $A_M$, and a complex scalar, $\Sigma$, responsible for the eventual $U(1)_{B-L}$ breaking living in the five-dimensional bulk. In string theory this set-up may be realized, for example, when the SM fields and the right-handed neutrinos live on a $(3 + 1)$-dimensional D-brane localized in the extra dimension, while the bulk fields arise from higher dimensional D-branes.

The potential for the scalar sector renormalizable in four dimensions is given by

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[4] One can make large $c_0$ natural by introducing a large number in the model, e.g. a multiplet with a large multiplicity.

| | $SU(3)_c$ | $SU(2)_L$ | $U(1)_Y$ | $U(1)_{B-L}$ |
|---|---|---|---|---|
| $q_L^i$ | **3** | **2** | $+1/6$ | $+1/3$ |
| $u_R^i$ | **3** | **1** | $+2/3$ | $+1/3$ |
| $d_R^i$ | **3** | **1** | $-1/3$ | $+1/3$ |
| $\ell_L^i$ | **1** | **2** | $-1/2$ | $-1$ |
| $\nu_R^i$ | **1** | **1** | $0$ | $-1$ |
| $e_R^i$ | **1** | **1** | $-1$ | $-1$ |
| $H$ | **1** | **2** | $-1/2$ | $0$ |
| $\Sigma$ | **1** | **1** | $0$ | $+2$ |

$$V(H, \Sigma) = \mu_H^2 H^\dagger H + \mu_\Sigma^2 \Sigma_0^* \Sigma_0 + \frac{\lambda_1}{2}\left(H^\dagger H\right)^2$$
$$+ \frac{\lambda_2}{2}\left(\Sigma_0^* \Sigma_0\right)^2 + \lambda_3 H^\dagger H \Sigma_0^* \Sigma_0. \tag{3.1}$$

Here, $\Sigma_0$ is the zero-mode of $\Sigma$ in the fifth direction. After spontaneous symmetry breaking, the scalar fields acquire vacuum expectation values (VEVs), and we can write

$$H = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_H + h \end{pmatrix}, \qquad \Sigma_0 = \frac{v_\Sigma + s}{\sqrt{2}}, \tag{3.2}$$

where $h$ and $s$ are excitations about the minimum, which is given by

$$\frac{v_H^2}{2} = \frac{-\mu_H^2 \lambda_2 + \mu_\Sigma^2 \lambda_3}{\lambda_1 \lambda_2 - \lambda_3^2}, \qquad \frac{v_\Sigma^2}{2} = \frac{-\mu_\Sigma^2 \lambda_1 + \mu_H^2 \lambda_3}{\lambda_1 \lambda_2 - \lambda_3^2}. \tag{3.3}$$

Note that the $W$ boson mass fixes $v_H = 246$ GeV. In terms of $h$ and $s$, the quadratic part of the potential is given by

$$V^{(2)} = \frac{1}{2}\eta^{\mathsf{T}} M_0^2 \eta, \quad \eta = \begin{pmatrix} h \\ s \end{pmatrix}, \tag{3.4}$$

where

$$M_0^2 = \begin{pmatrix} \lambda_1 v_H^2 & \lambda_3 v_H v_\Sigma \\ \lambda_3 v_H v_\Sigma & \lambda_2 v_\Sigma^2 \end{pmatrix}, \tag{3.5}$$

is the tree-level mass-squared matrix for $h$ and $s$, and we have used the minimization condition for the potential. Diagonalizing, the physical mass eigenstates are defined by

$$\begin{pmatrix} h \\ s \end{pmatrix} = \begin{pmatrix} \cos\alpha & \sin\alpha \\ -\sin\alpha & \cos\alpha \end{pmatrix}\begin{pmatrix} h_1 \\ h_2 \end{pmatrix}, \tag{3.6}$$

with the mixing angle given by

$$\tan 2\alpha = \frac{2\lambda_3 v_h v_s}{\lambda_2 v_s^2 - \lambda_1 v_h^2}. \tag{3.7}$$

The masses of the physical states are then given by

$$m_{h_{1,2}}^2 = \frac{1}{2}\left\{\lambda_1 v_H^2 + \lambda_2 v_\Sigma^2 \mp \sqrt{\left(\lambda_1 v_H^2 - \lambda_2 v_\Sigma^2\right)^2 + 4\lambda_3^2 v_H^2 v_\Sigma^2}\right\}. \tag{3.8}$$

For $|\lambda_3| \ll 1$ and $|v_H/v_\Sigma| \ll 1$, we can expand the square root and obtain

$$m_{h_{1,2}}^2 = \lambda_{1,2} v_{H,\Sigma}^2 \mp \frac{\lambda_3^2}{\lambda_2} v_H^2 + \lambda_3^2 v_H^2 \mathcal{O}\left(\frac{v_H^2}{v_\Sigma^2}\right). \tag{3.9}$$

Assuming no coupling between the Higgs $H$ and the scalar $\Sigma_0$ at tree level, the mixing term $H^\dagger H \Sigma_0^* \Sigma_0$ is induced at one-loop level [13] through interactions with neutrinos responsible for the seesaw mechanism [14–17]:
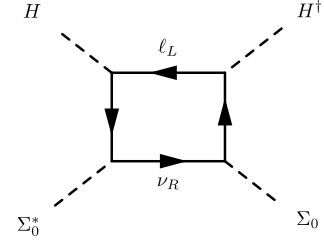


**Fig. 6.** One-loop diagram which contributes to the mixing term $H^\dagger H \Sigma_0^* \Sigma_0$ through the right-handed neutrinos.

$$\mathcal{L} \supset -Y_D^{ij} \overline{\nu_R^i} H^\dagger l_L^j - \frac{1}{2}Y_N^{ij} \overline{\nu_R^{ic}} \nu_R^j \Sigma_0 + \text{h.c.} \tag{3.10}$$

Fig. 6 displays the particular one-loop graph.[5] After $U(1)_{B-L}$ symmetry breaking, the mixing term contributes to the Higgs mass is estimated as

$$\delta m_H^2 \sim \frac{Y_D^2 Y_N^2}{(4\pi)^2}\frac{v_\Sigma^2}{2} \sim \frac{m_\nu M_N^3}{(4\pi)^2 v_H^2}, \tag{3.11}$$

where we have used the seesaw formula $m_\nu \sim Y_D^2 v_H^2 / M_N$ with $M_N = Y_N v_\Sigma / \sqrt{2}$ being the mass of the right-handed neutrino. Given the observation of the Higgs boson with mass 126 GeV at the LHC [18–20], we should have $\sqrt{|\delta m_H^2|} \lesssim 100$ GeV if naturalness is a criterion. Thus if we take $m_\nu \sim 0.1$ eV, we have $M_N \lesssim 10^7$ GeV from (3.11) and hence $v_\Sigma \lesssim 10^7/Y_N$ GeV. This translates to an upper bound on the mixing coupling

$$|\lambda_3| \sim \frac{Y_D^2 Y_N^2}{(4\pi)^2} \sim \frac{m_\nu M_N}{(4\pi)^2 v_H^2}Y_N^2 \lesssim 10^{-10}Y_N^2. \tag{3.12}$$

Assuming $Y_N \simeq \mathcal{O}(1)$, the mass of the physical $U(1)_{B-L}$ gauge boson is estimated as

$$m_{Z'} \sim g_4 v_\Sigma \lesssim \mathcal{O}\left(10^4\right) \text{ GeV}. \tag{3.13}$$

From collider experiments, one has $m_{Z'} \geqslant g_4 \times (6 \text{ TeV})$ for a $U(1)_{B-L}$ $Z'$ boson [21]. Since $g_4 \lesssim 10^{-3}$, there are no stringent bounds on $m_{Z'}$.

## 4. The inflaton decay

The coupling between the inflaton and the SM particles is crucial at the time of reheating. Let us first consider the following $\mathbb{Z}_2$ transformation:

$$x^5 \to -x^5, \qquad A_5 \to -A_5. \tag{4.1}$$

We choose the origin of the $x^5$ coordinate to be where the brane is localized. We assume there are no other fields with $\mathbb{Z}_2$-odd charges under (4.1) that are lighter than $A_5$. Then if this $\mathbb{Z}_2$ transformation is an exact symmetry, the inflaton is absolutely stable. This will be a problem, however, since then the Universe could not be heated to bring forth the standard hot big bang cosmology. We therefore introduce a five-dimensional Chern–Simons term, which breaks the $\mathbb{Z}_2$ symmetry:

$$S_{\text{CS}} = \frac{k}{48\pi^3}\int \mathcal{A}\mathcal{F}^2, \tag{4.2}$$

where $\mathcal{A} = \mathcal{A}_M dx^M$, $\mathcal{F} = \frac{1}{2}\mathcal{F}_{MN} dx^M dx^N$, and $k$ is some integer. Here, $\mathcal{A}_M$ is the $U(1)_{B-L}$ gauge field with mass dimension one, which is related to the canonically normalized fields by

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[5] Contributions from two-loop diagrams studied in [13] are suppressed in our model due to the smallness of the $U(1)_{B-L}$ gauge coupling.

$$A_M^{(5)} = \frac{1}{g_5} \mathcal{A}_M, \tag{4.3}$$

$$A_\mu = \frac{1}{g_4} \mathcal{A}_{\mu 0}, \qquad \phi = \frac{1}{g_4} \mathcal{A}_{50}, \tag{4.4}$$

where $A_M^{(5)}$ is the $U(1)_{B-L}$ gauge field canonically normalized in five dimensions, $A_\mu$ that in four dimensions, and $\mathcal{A}_{M0}$ the zero-mode of $\mathcal{A}_M$ in five dimensions.

The four-dimensional interaction of the zero-modes following from (4.2) is

$$\frac{k}{48\pi^3} \int d^4x \, (2\pi L) \frac{3}{4} \epsilon^{\mu\nu\rho\sigma} \mathcal{A}_{50} \mathcal{F}_{\mu\nu 0} \mathcal{F}_{\rho\sigma 0}$$
$$= \frac{k}{16\pi^2} \int d^4x \, \frac{\phi}{2\pi f} \mathcal{F}_{\mu\nu 0} \tilde{\mathcal{F}}_0^{\mu\nu}$$
$$= g_4^2 \frac{k}{16\pi^2} \int d^4x \, \frac{\phi}{2\pi f} F_{\mu\nu} \tilde{F}^{\mu\nu}. \tag{4.5}$$

Here the subscript 0 denotes that they are (made from) zero-modes in the fifth direction.

The coupling (4.5) gives the dominant contribution to the decay width at the tree level:

$$\Gamma_{\phi \to AA} \simeq \frac{g_4^4}{16\pi} \left( \frac{k}{32\pi^3} \right)^2 \frac{m_\phi^3}{f^2}, \tag{4.6}$$

where $m_\phi$ is the mass of the inflaton. As we have seen, $c_0$ is determined by (2.18) once $f$ and $g_4$ are given. This then determines $m_\phi$:

$$\frac{m_\phi^2}{2} = \frac{V_0}{4f^2} = \frac{1}{4f^2} \frac{c_0}{\pi^2} \frac{1}{(2\pi L_5)^4} = \frac{c_0 g_4^4 f^2}{4\pi^2}. \tag{4.7}$$

The $U(1)_{B-L}$ gauge bosons decay to SM particles via the minimal couplings. As this proceeds much faster than the inflaton decay, the reheating temperature is governed by the inflaton decay width (4.6). It is estimated as

$$T_R = \left( \frac{90}{\pi^2 g_\star(T_R)} \right)^{1/4} \sqrt{\Gamma M_P}$$
$$\simeq \left( \frac{90}{g_\star(T_R)} \right)^{1/4} \frac{g_4^2}{4\pi} \frac{|k|}{32\pi^3} \sqrt{\frac{m_\phi^3 M_P}{f^2}}$$
$$\simeq |k| \times 1 \text{ GeV}, \tag{4.8}$$

where in the last line, we have used the preferred values $f \simeq 5M_P$ and $\ell_5 \simeq 5$, which gives $m_\phi \simeq 10^{13}$ GeV. The factor $g_\star(T)$ is the effective relativistic degrees of freedom at temperature $T$. For $T_R \simeq 1 \sim 10$ GeV, $g_\star(T_R) \simeq 60 \sim 90$. From (4.8), the reheating temperature is much smaller than the $U(1)_{B-L}$ breaking scale given by $v_\Sigma \simeq \mathcal{O}(10^7)$ GeV, when $k$ is $\mathcal{O}(1-10)$. Comparing (4.8) with the standard estimate of the number of e-folds [22]:

$$N_* \simeq 49 + \frac{2}{3} \ln \left( \frac{\rho_*^{1/4}}{10^{16} \text{ GeV}} \right) + \frac{1}{3} \ln \left( \frac{T_R}{1 \text{ GeV}} \right), \tag{4.9}$$

we observe that $N_* \simeq 50$ is natural in our model, as advertised earlier.

## 5. Summary and discussions

In this Letter, we have studied the interrelation between cosmology and particle physics in $U(1)_{B-L}$ extra-natural inflation with a gauged $U(1)_{B-L}$ extension of the SM localized on a brane. The cosmological observation constrains the value of the $U(1)_{B-L}$

gauge coupling to $g_4 \lesssim 10^{-3}$, which in turn constrains the particle physics scenario at high energy assuming naturalness. On the other hand, with SM particles localized on a brane, allowed interaction between the inflaton and the SM particles are restricted. Together with the value of $g_4$, the decay width of the inflaton and the reheating temperature are determined.

By tuning of a few parameters or with some slight extension, our model may also be able to explain other cosmological observations such as the Baryon number asymmetry of the Universe and the dark matter abundance. Indeed, the right-handed neutrinos could play a role in the former through leptogenesis,[6] and they are also dark matter candidates. Another possible dark matter candidate, which may be included in our model, is a light scalar field odd under the reflection of the extra dimension (4.1). These merit further investigations.

Our main purpose in this Letter is to present an example in which the relation between the BSM physics and the inflation physics are specified, and theoretical and observational constraints on one side constrains the other. We discussed one example here, but there can be several other possibilities, even within gauged $U(1)_{B-L}$ extensions of the SM. For instance, one may put some of the SM fields in the bulk. It will be interesting to explore those related scenarios.

## Acknowledgements

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