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$U(1)_{B-L}$ symmetry restoration and effective neutrino species

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ABSTRACT

The $U(1)_{B-L}$ symmetry could be restored during inflation, since the BICEP2 results suggest a GUT-scale inflation with the Hubble parameter, $H_{\rm inf} \simeq 10^{14}$ GeV, close to the $U(1)_{B-L}$ breaking scale. We consider a scenario in which the $B-L$ Higgs field dominates the Universe after inflation, and mainly decays into the $U(1)_{B-L}$ gauge bosons, whose subsequent decays reheat the Universe. Interestingly, if one of the right-handed neutrinos is extremely light and behaves as dark radiation or hot dark matter, its abundance is determined by the $B-L$ charge assignment and the relativistic degree of freedom in plasma. We find that $\Delta N_{\rm eff}$ takes discrete values between 0.188 and 0.220 in the standard model plus three right-handed neutrinos, depending on whether the decay into heavier right-handed neutrinos is kinematically accessible or not. In the fiveness $U(1)_{\mathbf{5}}$ case, we find that $\Delta N_{\rm eff}$ takes discrete values between 0.313 and 0.423. The tension between BICEP2 and *Planck* can be partially relaxed by dark radiation.

1. Introduction

The BICEP2 experiment detected the primordial B-mode polarization of cosmic microwave background (CMB) with a high significance [1]. This could be due to tensor mode perturbations generated during inflation, and if correct, it suggests a rather high inflation scale:

$$H_{\rm inf} \simeq 1.0 \times 10^{14}\ {\rm GeV} \left(\frac{r}{0.16}\right)^{\frac{1}{2}}, \tag{1}$$

$$r = 0.20^{+0.07}_{-0.05}\ (68\%{\rm CL}). \tag{2}$$

For such high-scale inflation, various symmetries may be restored during inflation. Also, some of the symmetries broken during inflation can be restored after inflation, if the reheating temperature is sufficiently high. In this Letter we revisit cosmological implications of such symmetry restoration and its subsequent breaking.

Among various symmetries, we consider an extra U(1) gauge symmetry, which is assumed to be restored during (or after) inflation and become spontaneously broken sometime after inflation.[1] We mainly focus on the $U(1)_{B-L}$ symmetry as such, since it is a plausible extension of the standard model (SM) motivated by grand unification theory (GUT) as well as the charge quantization argument in the presence of three right-handed neutrinos. If exists, the $U(1)_{B-L}$ symmetry must be spontaneously broken in the present vacuum, and the breaking scale is expected to be of order 10^{13-16} GeV based on the measured neutrino mass squared differences and the seesaw mechanism [6].

One of the straightforward consequences of the $U(1)_{B-L}$ breaking after inflation is the production of cosmic strings, which can be searched for by the CMB observations [7] as well as pulsar timing measurements [8]. Another is the dynamics of the $B-L$ Higgs field during the phase transition. In particular, as studied in Refs. [9,10], there may be a phase during which the $B-L$ Higgs field, being trapped at the origin, induces a mini-inflation or thermal inflation [11–14]. Then the Universe after mini-inflation will be dominated by the $B-L$ Higgs, whose decay reheats the SM sector. This scenario has an advantage that the huge entropy produced by the $B-L$ Higgs decay relaxes the overproduction of unwanted relics such as gravitinos from the inflaton decay [15–18].

Alternatively, it is possible that the $B-L$ Higgs field plays a role of the inflaton. For instance, a quadratic chaotic inflation can be realized if its kinetic term is modified at large field values, as in the running kinetic inflation [19–22]. In this case, the $B-L$ Higgs field necessarily dominates the Universe after inflation.

If kinematically allowed, the $B-L$ Higgs field can mainly decay into the $B-L$ gauge bosons. This is the case if the right-handed neutrinos are either heavier than a half of the $B-L$ Higgs boson mass or much lighter. Then, the Universe will be reheated

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[1] The implications of the BICEP2 results for a global U(1) Peccei–Quinn symmetry [2] and the axion cold dark matter has been discussed in Refs. [3–5].

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by decays of the $B-L$ gauge bosons. Interestingly, the branching fractions of various decay modes are then determined solely by the $B-L$ charge assignment. If all the decay products enter thermal equilibrium, the initial branching ratios will be soon forgotten without any consequences in the low energy. Some of the decay products, however, may stay out-of-equilibrium until today, retaining the valuable information of the beginning of the radiation dominated Universe. One plausible candidate in the minimal extension of SM is the right-handed neutrinos. Indeed, if the effective mass of the lightest right-handed neutrino is of order keV, it can be warm dark matter [23–25],[2] and if it is much lighter, it can contribute to the effective neutrino species as dark radiation or hot dark matter. We consider the latter possibility in this letter. The presence of dark radiation or hot dark matter can relax the tension between BICEP2 and Planck [28].

In the next section we will first discuss the $B-L$ breaking scale suggested by the seesaw formula, and study the cosmological evolution of the $B-L$ Higgs field. Then we estimate the contribution of the lightest right-handed neutrino to the effective neutrino species in the case of $\mathrm{U}(1)_{B-L}$ and the so-called fiveness $\mathrm{U}(1)_{\mathbf{5}}$. The last section is devoted for discussion and conclusions.

2. $B-L$ Higgs cosmology and effective neutrino species

2.1. Seesaw mechanism and $\mathrm{U}(1)_{B-L}$ breaking scale

First let us review the seesaw mechanism to estimate the typical breaking scale of the $\mathrm{U}(1)_{B-L}$ symmetry. We extend the SM by adding three right-handed neutrinos and consider the interactions,

$$\mathcal{L} = i\bar{N}_I\gamma^\mu\partial_\mu N_I - \left(\lambda_{I\alpha}\bar{N}_I L_\alpha H + \frac{1}{2}\kappa_I \Phi \bar{N_I^c} N_I + \text{h.c.}\right), \tag{3}$$

where N_I, L_α, H and Φ are the right-handed neutrino, lepton doublet and Higgs scalar, the $B-L$ Higgs scalar, respectively, I denotes the generation of the right-handed neutrinos, and α runs over the lepton flavor, e, μ and τ. The sum over repeated indices is understood. After the spontaneous breakdown of the $\mathrm{U}(1)_{B-L}$ gauge symmetry, the right-handed neutrinos acquire a mass,

$$M_I = \kappa_I\langle\Phi\rangle. \tag{4}$$

Here we adopt a basis in which the right-handed neutrinos are mass eigenstates with $M_1 \le M_2 \le M_3$. The seesaw formula for the light neutrino mass is obtained by integrating out the heavy right-handed neutrinos:

$$(m_\nu)_{\alpha\beta} = \lambda_{\alpha I}\lambda_{I\beta}\frac{v^2}{M_I}, \tag{5}$$

where $v \equiv \langle H^0\rangle \simeq 174$ GeV is the vacuum expectation value (VEV) of the Higgs field. As a typical neutrino mass scale, we adopt the mass squared difference measured by the atmospheric neutrino oscillation experiments, $m_\nu \simeq 0.05$ eV. Then the $B-L$ breaking scale inferred from the seesaw formula ranges as

$$\langle\Phi\rangle \approx 10^{13}\text{–}10^{16}\ \text{GeV} \tag{6}$$

for $\lambda_{\alpha I} = \mathcal{O}(0.1-1)$ and $\kappa_I = \mathcal{O}(0.1-1)$. Since the $B-L$ breaking scale is close to the Hubble parameter during inflation suggested by the BICEP2 results, it is possible the $\mathrm{U}(1)_{B-L}$ symmetry is restored during inflation.[3] This is especially the case if the breaking scale is close to the lower end of the above range (6).

[2] See Refs. [26,27] for the implications for the 3.5 keV X-ray line.

[3] For instance, a non-minimal coupling to the gravity, $\xi|\Phi|^2R$, can stabilize the origin of Φ for a certain value of ξ.

Lastly let us note that some of the right-handed neutrinos can have a mass much smaller than the typical $B-L$ breaking scale. In fact, it is known that the above mentioned feature of the seesaw formula can be preserved even for a split mass spectrum of the right-handed neutrinos in the simple Froggatt–Nielsen model [29] or the split seesaw mechanism [9]. Also, it is possible to make the lightest one, N_1, extremely light so that it does not contribute to the light neutrino mass, in the split flavor model [10,26].[4] It is of course possible to make N_1 massless by imposing a certain flavor symmetry on only N_1. Later we shall consider a case in which N_1 is so light that it behaves as dark radiation or hot dark matter.

2.2. $B-L$ Higgs-dominated Universe

Let us here briefly discuss two scenarios in which the $B-L$ Higgs field dominates the energy density of the Universe after inflation. In the first scenario we assume that $\mathrm{U}(1)_{B-L}$ symmetry is restored during inflation, and the $B-L$ Higgs, being trapped at the origin, drives a mini-thermal inflation. In the second scenario, we consider a case in which the $B-L$ Higgs field plays a role of the inflaton rolling down the potential from large field values. This is possible if the kinetic term runs at large field values [19–22].

The potential for the $B-L$ Higgs field Φ is given by

$$V(\phi) = -\frac{1}{2}\mu^2\phi^2 + \frac{\lambda}{4}\phi^4, \tag{7}$$

where we have defined $\phi = \sqrt{2}|\Phi|$. In the present vacuum ϕ develops a vacuum expectation value (VEV) as

$$\langle\phi\rangle = \frac{\mu}{\lambda}, \tag{8}$$

which is considered to be within the range of (6). The mass of the $B-L$ Higgs boson at the low-energy minimum is $m_\phi = \sqrt{2}\mu$. As a reference value, we take $\langle\phi\rangle \approx 10^{13}$ GeV. Then, even if the $\mathrm{U}(1)_{B-L}$ is broken during inflation, it can be restored after inflation, if the reheating temperature is sufficiently high, $T_R \gtrsim 10^{13}$ GeV.

Let us suppose that the $B-L$ Higgs field is trapped at the origin after inflation and therefore $\mathrm{U}(1)_{B-L}$ is restored. Taking account of the thermal effects,[5] the potential around the origin can be written as

$$V \approx V_0 + \frac{1}{2}\left(c_g g_{B-L}^2 + c_\lambda\lambda + c_\kappa\kappa_3^2\right)T^2\phi^2 - \frac{1}{2}\mu^2\phi^2 + \cdots, \tag{9}$$

where $V_0 = \mu^4/4\lambda$, c_g, c_λ and c_κ are numerical coefficients of order $\mathcal{O}(0.1)$, g_{B-L} denotes the gauge coupling of $\mathrm{U}(1)_{B-L}$, κ_3 denotes the coupling of the $B-L$ Higgs to the heaviest right-handed neutrino, and T is the temperate of the background thermal plasma. For sufficiently high temperature, ϕ is stabilized at the origin. The critical temperature at which the origin becomes unstable is given by

$$T_c \simeq \frac{\mu}{\sqrt{c_g g_{B-L}^2 + c_\lambda\lambda + c_\kappa\kappa_3^2}}. \tag{10}$$

The condition for the $B-L$ Higgs to dominate the Universe at the critical temperature reads

[4] We can achieve both sufficiently small mass and mixing simultaneously so that production from the Dodelson–Widrow mechanism [30] is negligible. In the split flavor mechanism, the breaking of flavor symmetry is tied to the breaking of $B-L$ symmetry. The spontaneous breakdown of $\mathrm{U}(1)_{B-L}$ may lead to the formation of domain walls, which however can be removed if the flavor symmetry is only approximate.

[5] Here we assume that the inflaton decays into the SM particles so that there is dilute hot plasma.

$$V_0 \gtrsim T_c^4 \quad \Leftrightarrow \quad \left(c_g g_{B-L}^2 + c_\lambda \lambda + c_\kappa \kappa_3^2\right)^2 \gtrsim \lambda. \tag{11}$$

This can be satisfied for $\lambda = \mathcal{O}(1)$. Even for small λ, the condition can be met for $\kappa_3 = \mathcal{O}(1)$. Note that a large κ_3 is needed in this case since we are interested in the case where the $B-L$ Higgs decays mainly into the $B-L$ gauge bosons, which requires $g_{B-L}^2 \lesssim \lambda$.

Once the $B-L$ Higgs field dominates the Universe, those particles produced before the domination will be diluted by the subsequent decay of the $B-L$ Higgs. In particular, we assume that the thermal population of N_1 formed before the domination gives only negligible contributions to the final abundance in the following.

Alternatively we can consider a case in which the $B-L$ Higgs field plays the role of the inflaton. This is possible if the kinetic term depends on the $B-L$ Higgs field itself as [19–22]

$$\mathcal{L}_K = \frac{1}{2}\left(1 + \xi\phi^2\right)(\partial\phi)^2, \tag{12}$$

where $\xi \gtrsim 1/M_p^2$ is the coupling constant, and $M_p \simeq 2.4 \times 10^{18}$ GeV. At sufficiently large field values, $\phi \gtrsim 1/\sqrt{\xi}$, the canonically normalized field is given by $\hat{\phi} \sim \sqrt{\xi}\phi^2$, and therefore the quartic potential for ϕ turns into the mass term for $\hat{\phi}$ with the mass $m_{\hat{\phi}}^2 \sim \lambda/\xi$. Thus the quadratic chaotic inflation model is realized by the $B-L$ Higgs field with the running kinetic term, which is consistent with the BICEP2 results (1). In this case, the Universe after inflation is naturally dominated by the $B-L$ Higgs field.

In addition to the above scenarios, there are various possibilities to realize the $B-L$ Higgs-dominated Universe. For instance, one may consider a short duration of the hybrid inflation [31] with the waterfall field being identified with the $B-L$ Higgs field. In contrast to the usual hybrid inflation, the $B-L$ Higgs field can have a mass comparable to the $B-L$ breaking scale.

2.3. Decays of $B-L$ Higgs

Here let us study the decays of the $B-L$ Higgs ϕ. The decay rate for $\phi \to 2N_I$ and $\phi \to 2A_\mu$ are given as

$$\Gamma_{\phi\to 2N_I} = \frac{1}{8\pi}\kappa_I^2 m_\phi \left(1 - \frac{4M_I^2}{m_\phi^2}\right)^{3/2}, \tag{13}$$

$$\Gamma_{\phi\to 2A_\mu} \approx \frac{g_{B-L}^2}{128\pi}\frac{m_\phi^3}{m_A^2} \tag{14}$$

where $m_A = g_{B-L}\langle\phi\rangle$ is the $B-L$ gauge boson mass, and we have approximated $m_\phi \gtrsim 2m_A$.

We would like to consider a situation where the $B-L$ Higgs mainly decays into the $B-L$ gauge bosons. To this end, we require

$$\Gamma_{\phi\to 2N_I} \ll \Gamma_{\phi\to 2A_\mu}. \tag{15}$$

We are interested in a case where N_1 is much lighter than the other two, i.e., $M_1 \ll M_2, M_3$, and so, practically the decay into N_1 is negligible. Let us focus on the heaviest right-handed neutrino N_3. The same analysis also holds for N_2. If κ_3 is of order unity and $\lambda \lesssim \mathcal{O}(0.1)$, the decay into a pair of N_3 can be kinematically forbidden. In this case (15) is automatically satisfied. On the other hand, if it is kinematically accessible, the above condition places an upper bound on κ_3,

$$\kappa_3 \ll \frac{g_{B-L}}{4}\left(\frac{m_\phi}{m_{A_\mu}}\right) = \frac{\lambda}{2\sqrt{2}}. \tag{16}$$

Thus, as long as $\lambda = \mathcal{O}(1)$, the above condition is satisfied if κ_3 is smaller than $\mathcal{O}(10^{-2})$. A similar argument holds for N_2.

When the $B-L$ Higgs starts to oscillate from large field values, it efficiently dissipates its energy into thermal plasma, producing the $B-L$ gauge bosons as well as the right-handed neutrinos [32]. If κ_I is sufficiently small, we can suppress the production of the right-handed neutrinos with respect to that of the $B-L$ gauge bosons. Although it depends on the details of the thermalization processes, it is possible that the main reheating process is through the perturbative decays of the $B-L$ gauge bosons, which are non-perturbatively produced by the inflaton dynamics. This is the case if the relevant dissipation proceeds like the instant preheating [33]. Then our scenario is applicable to this case as well.

2.4. Effective neutrino species

2.4.1. $U(1)_{B-L}$ symmetry

The lightest right-handed neutrino produced by decays of the $B-L$ gauge bosons will increase the effective number of neutrino species ($N_{\rm eff}$) by the amount [34,35]

$$\Delta N_{\rm eff} = \left.\frac{\rho_{N_1}}{\rho_\nu}\right|_{\nu\ \rm decouple} = \frac{43}{7}\frac{B_1}{1-B_1}\left(\frac{43/4}{g_*(T_d)}\right)^{1/3}, \tag{17}$$

where B_1 is the branching fraction of the $B-L$ gauge bosons to a pair of N_1, and $g_*(T_d)$ counts the relativistic degrees of freedom in thermal plasma at the decay of the $B-L$ gauge bosons. In deriving the above expression, we have used the fact that the entropy in the comoving volume is conserved.

We are interested in the following three cases: (i) $M_2 \leq M_3 \ll m_\phi$; (ii) $M_2 \ll m_\phi < 2M_3$; (iii) $m_\phi < 2M_2 \leq 2M_3$. In these cases, the $B-L$ Higgs mainly decays into the $B-L$ gauge bosons. Then branching fraction into N_1 is given by $B_1 = 1/16, 1/15$ and $1/14$ for the cases (i), (ii), and (iii), respectively. This leads to the robust prediction of $\Delta N_{\rm eff}$ as

$$\Delta N_{\rm eff} \simeq \begin{cases} 0.188 & \text{case (i)} \\ 0.203 & \text{case (ii)} \\ 0.220 & \text{case (iii)}, \end{cases} \tag{18}$$

where we have assumed that the decay products (including the heavy right-handed neutrinos) enter thermal equilibrium. This assumption is used to evaluate $g_*(T_d)$, to which our results are not sensitive.

2.4.2. Fiveness $U(1)_{\mathbf{5}}$ symmetry

We can also consider a certain mixture of $\mathrm{U}(1)_{B-L}$ and $\mathrm{U}(1)_Y$, the so-called fiveness $\mathrm{U}(1)_{\mathbf{5}}$, based on a GUT model with a symmetry breaking pattern $\mathrm{SO}(10) \to \mathrm{SU}(5) \times \mathrm{U}(1)_{\mathbf{5}}$. The charges of the $B-L$, fiveness and hyper charge are related as [36]

$$B - L = \frac{1}{5}Y_5 + \frac{4}{5}Y, \tag{19}$$

that is, sterile neutrinos transform as $(\mathbf{1}, +5)$.

In this case, there are Higgs fields, Φ_5 and $\Phi_{\bar{5}}$, which transform as $(\mathbf{5}, -2)$ and $(\bar{\mathbf{5}}, 2)$. These Higgs fields contain colored Higgs as well as two Higgs doublets, and we assume that the colored Higgs are heavier than the $B-L$ Higgs boson. The SM Higgs doublet is given by a certain combination of the two Higgs doublets. In addition to the cases (i)–(iii) considered before, there are two cases we can consider; case (A): the two Higgs doublets are lighter than $m_\phi/2$; case (B): one of the two Higgs doublets is heavier than $m_\phi/2$.

In the case (i) with $M_2 \leq M_3 \ll m_\phi$, the branching fraction of the $B-L$ gauge boson into the lightest right-hand neutrinos is given by $B_1 = 25/248$ and $24/244$ for the cases (A) and (B), respectively. Here we have taken into a fact that the partial decay rate of the $B-L$ gauge boson into scalars is half of that into fermions with the same charge. Then we can estimate $\Delta N_{\rm eff}$ as

$$\Delta N_{\rm eff} \simeq \begin{cases} 0.313 & \text{case (A)} \\ 0.323 & \text{case (B)} \end{cases} \tag{20}$$

Similarly, in the case (ii) with $M_2 \ll m_\phi < 2M_3$, we obtain $B_1 = 25/223$ and $25/219$ for the cases (A) and (B), respectively, and $\Delta N_{\rm eff}$ is given by

$$\Delta N_{\rm eff} \simeq \begin{cases} 0.355 & \text{case (A)} \\ 0.366 & \text{case (B)} \end{cases} \tag{21}$$

Lastly, in the case (iii) with $m_\phi < 2M_2 \le 2M_3$, we obtain $B_1 = 25/198$ and $25/194$ for the cases (A) and (B), respectively, and $\Delta N_{\rm eff}$ is given by

$$\Delta N_{\rm eff} \simeq \begin{cases} 0.408 & \text{case (A)} \\ 0.423 & \text{case (B)} \end{cases} \tag{22}$$

Thus, the effective neutrino species tends to be larger than the case of $\mathrm{U}(1)_{B-L}$.

3. Discussion

We have so far considered the case in which the $B-L$ Higgs field dominates the Universe and mainly decays into the $B-L$ gauge bosons, in order to ensure that the branching fractions of various decay processes are simply determined by the $B-L$ charge assignment. There are other possibilities to realize the robust prediction of $\Delta N_{\rm eff}$. For instance, one can consider a hidden U(1) gauge symmetry, which has a kinetic mixing with $\mathrm{U}(1)_{B-L}$. Assuming that there are no matter fields charged under the hidden U(1) symmetry, the hidden gauge boson decays into the SM particles through the kinetic mixing with $\mathrm{U}(1)_{B-L}$ [37]. In this case, the branching fractions of the decay processes are similarly determined by the $B-L$ charge assignment. Instead of hidden gauge bosons, one can also consider hidden gaugino as well. In order for the hidden gauge bosons (or hidden gauginos) to dominate the Universe, one may consider that the inflation takes place in the hidden sector. For instance, one may identify the hidden Higgs field with the inflaton. Then most of the above arguments can be applied to the hidden Higgs dynamics.

The baryon asymmetry can be created through leptogenesis [38]. In the present scenario there are two heavy right-handed neutrinos, and the decay of N_2 can generate the right amount of the baryon asymmetry for $M_2 \gtrsim 10^{11}$ GeV [39,40]. Taking $\langle\phi\rangle = \mathcal{O}(10^{13-14})$ GeV, it is possible to suppress the direct decay of the $B-L$ Higgs into a pair of N_2 so that our results about $\Delta N_{\rm eff}$ remain intact.

So far we have assumed that the direct decay of the $B-L$ Higgs into N_2 and N_3 are suppressed. If the partial decay rate into N_2 or N_3 becomes comparable to or even larger than that into $B-L$ gauge bosons, the abundance of extra neutrino species is suppressed. In this sense our results on $\Delta N_{\rm eff}$ can be thought of as the upper bound in a scenario where the $B-L$ Higgs dominates the Universe and the lightest right-handed neutrinos behaves as dark radiation or hot dark matter.

We have taken up two examples, $\mathrm{U}(1)_{B-L}$ and $\mathrm{U}(1)_{\mathbf{5}}$, to show that the additional effective neutrino specifies can be fixed by the charge assignment and the particle contents. Therefore, these predictions on $\Delta N_{\rm eff}$ are robust, and can be tested in future CMB experiments, which will achieve $\sigma(N_{\rm eff}) \simeq 0.02$ [41]. There are two ways to extend our results. One is to enlarge the particle content. For instance, it was discussed in Ref. [42] how one can add chiral fermions charged under the $\mathrm{U}(1)_{B-L}$ satisfying the anomaly cancellation conditions. If some of the extra fermions are sufficiently light, we can increase $\Delta N_{\rm eff}$ in a similar manner. Alternatively, we may apply our idea to different gauge symmetry. In particular, it is straightforward to consider another possible U(1) extensions based on the GUT group with a higher rank, such as E_6 [43]. In this case we may have to introduce a flavor symmetry on the extra fermions to ensure their light mass.

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