Higgs vacuum stability and inflationary dynamics after BICEP2 and PLANCK dust polarisation data

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Abstract. If the recent detection of $B$-mode polarization of the Cosmic Microwave Background by BICEP2 observations, withstand the test of time after the release of recent PLANCK dust polarisation data, then it would surprisingly put the inflationary scale near Grand Unification scale if one considers single-field inflationary models. On the other hand, Large Hadron Collider has observed the elusive Higgs particle whose presently observed mass can lead to electroweak vacuum instability at high scale ($\sim \mathcal{O}(10^{10})$ GeV). In this article, we seek for a simple particle physics model which can simultaneously keep the vacuum of the theory stable and yield high-scale inflation successfully. To serve our purpose, we extend the Standard Model of particle physics with a $U(1)_{B-L}$ gauged symmetry which spontaneously breaks down just above the inflationary scale. Such a scenario provides a constrained parameter space where both the issues of vacuum stability and high-scale inflation can be successfully accommodated. The threshold effect on the Higgs quartic coupling due to the presence of the heavy inflaton field plays an important role in keeping the electroweak vacuum stable. Furthermore, this scenario is also capable of reheating the universe at the end of inflation. Though the issues of Dark Matter and Dark Energy, which dominate the late-time evolution of our universe, cannot be addressed within this framework, this model successfully describes the early universe dynamics according to the Big Bang model.

Keywords: inflation, particle physics - cosmology connection, physics of the early universe, neutrino properties

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1 Introduction

The recent observations, in the fields of both particle physics and cosmology, are immensely exciting as they have led to new discoveries in their respective fields. On the one hand, in particle physics the discovery of the elusive Higgs boson at the Large Hadron Collider (LHC) [1, 2] completes the picture of the Standard Model (SM) of particle physics. On the other hand, in the field of cosmology the recent observation of $B$-mode polarization of the Cosmic Microwave Background Radiation (CMBR) by BICEP2 [3], if confirmed, would reveal the picture of the primordial universe through the effect of tensor perturbations, which had defied detection so far.

We know new discoveries often lead to new conundrums and these observations are no exceptions to that. They have also set profound challenges in the field of particle physics. The observed mass of the Higgs boson ($\sim 122–127$ GeV) at the LHC [1, 2] may lead to an unstable electroweak vacuum in higher energy scales. This is because the scalar quartic coupling eventually becomes negative at very high scale leaving the scalar potential unbounded from below. Hence it is of importance to address the issue of stability of the SM vacuum [4] which has been analyzed in [5–15]. It has been noted, that it is indeed a difficult task to maintain the stability of the SM vacuum till the Planck scale [13, 14, 16], which indicates towards the onset of new physics before the electroweak vacuum becomes unstable [16–22].

Detection of $B$-mode polarization by BICEP2 experiment [3], if confirmed, would put inflationary paradigm [23, 24] on stronger footing than ever as only the primordial tensor modes produced during inflation are capable of imprinting $B$-mode polarization on the CMBR. In the simplest inflationary picture, the primordial universe during inflation is dominated by the potential energy of a slowly rolling scalar field, called inflaton, whose particle physics origin is yet to be known. Non-detection of isocurvature modes and primordial non-Gaussianity by the recent PLANCK observation [25] favour this simplest single-field model of inflation. The
aim of the BICEP2 experiment was to measure the ratio of the amplitudes of the primordial tensor and the scalar perturbations, known as the tensor-to-scalar ratio denoted by $r$. PLANCK put an upper limit on tensor-to-scalar ratio as $r < 0.11$, whereas BICEP2 claimed to have measured $r$ at large angular scales as

$$r = 0.2^{+0.07}_{-0.05},$$

where $r = 0$ is ruled out at $7\sigma$ confidence level. Such dispute in observations can be overcome if in future the running of the scalar spectral index $(dn_s/d\ln k)$ turns out to be large and negative [3]. Recently after the release of PLANCK’s polarised dust data [26] the BICEP2’s detection of $B$-mode polarisation has been put under severe scrutiny. Despite such doubt, a high-enough detection of $r$ has profound importance in cosmology for many reasons. First of all, such a high value of $r$, as has been claimed to have been observed by BICEP2, indicates a large-field inflationary scenario such as chaotic inflation [24] and favours quadratic as well as quartic potentials of the inflaton field [27]. Secondly, the observed value of $r$ by BICEP2 sets the scale of inflation surprisingly close to the Grand Unification (GUT) scale ($\sim 10^{16}$ GeV). Furthermore, as one requires the dynamics of a scalar field (preferably a fundamental scalar field) for inflation and the only known candidate in nature so far is the Higgs field, it is most economic to let the Higgs play the role of inflaton. Such Higgs inflationary scenario was first proposed in ref. [28] where the Higgs field strongly couples to the Ricci scalar. After the announcement of BICEP2 results, it has been argued that detection of such high $r$ has made it difficult for the SM Higgs boson to be the inflaton during inflationary era as these scenarios generally produce low tensor mode amplitude ($r \sim \mathcal{O}(10^{-2})$) [29, 30].

There are counter arguments, based on the roles of the top and the Higgs mass during the inflationary period, to rescue the Higgs inflation scenario [31–33]. To revive such possibilities few attempts have been made by incorporating non-minimal coupling between the Higgs kinetic term and the Higgs field [34], or the Einstein tensor [35] or both [36], or by adding a cosmological constant [37] or by introducing extra singlet scalar dark matter [38]. From the discussion so far we can infer that:

- The electroweak vacuum of the SM is not stable all the way up to the Planck scale if one does not invoke new physics in between the electroweak scale and the Planck scale.
- In a generic Higgs inflationary scenario the SM Higgs field, not being able to produce high enough $r$ according to BICEP2 observations, is not favoured as inflaton. Thus one might take an alternative path to identify some other scalar field as inflaton, while keeping in mind that the SM Higgs essentially plays no dynamical role during inflation in order to yield a single-field inflationary epoch favoured by PLANCK.

If one assumes that no other new physics sets in till inflationary scale, which now according to BICEP2 observation takes place at the GUT scale, then one needs to extend the SM scenario to the inflationary scale where introducing new physics one can account for the inflaton. It may happen that this new physics input in the GUT scale can make the Higgs vacuum stable up to the Planck scale. In this article we will deal with a minimal extension of the SM which can explain the vacuum stability of the SM as well as the inflationary dynamics in the light of PLANCK and BICEP2 results. Furthermore, it is well known that

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1It has recently been observed that high-scale inflation, indicated by BICEP2 observations, would destabilize the SM Higgs vacuum during inflation and one thus requires to invoke small coupling between the inflaton and the SM Higgs field in order to avoid such catastrophe [39].
in non-supersymmetric theories the quadratic divergence of the Higgs mass remains an open problem and few attempts have been made to deal with such difficulties [40]. In this article we would also try to address whether the Higgs mass can be kept light at the inflationary scale so that reheating at the end of inflation can take place through the channel of inflaton decay into SM Higgs field. There are few other very important outstanding issues with the present Standard Models of particle physics and cosmology, such as the origin of Dark Energy and Dark Matter, which are beyond the scope of the present article.

To serve our purpose, we extend the SM by an additional $U(1)_{B-L}$ gauge symmetry. Phenomenological aspects and the issues related to the vacuum stability of this $U(1)_{B-L}$ extended SM have been extensively analyzed in [20, 21, 41–46]. Such extension of SM has also been discussed previously in several cosmological contexts such as in inflationary scenarios [47, 48], to explain the origin of dark matter [49–51], baryogenesis and leptogenesis [52, 53] and production of gravitational waves [54]. Here we consider the spontaneous breaking of the $U(1)_{B-L}$ symmetry and the electro-weak symmetry to take place at very different energy scales. While the former takes place above the inflationary scale ($\sim 10^{16}$ GeV) and the real part of the scalar of this symmetry group plays the role of inflaton, the SM electroweak symmetry breaking takes place around 246 GeV. Also, to our advantage, couplings between the scalar of the $U(1)_{B-L}$ and the SM particles help reheat the universe at the end of inflation.

We have organized the rest of the article as follows. In section 2 we explain our particle physics model in the cosmological backdrop. In section 3 the constraints on model parameters from vacuum stability and inflationary dynamics are presented and the last section summarizes the main results obtained in the present work.

2 The model

In this section we discuss the particle physics model, the SM extended by an Abelian $U(1)_{B-L}$ group, which we will pursue to unify the concepts of inflation and electroweak vacuum stability.

2.1 Scalar sector

The $U(1)_{B-L}$ gauged extended SM, where the full symmetry group is depicted as

$$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_{B-L},$$

contains three extra right handed neutrinos to cancel all the gauge as well as gravitational anomalies, one extra gauge boson and one extra heavy scalar field ($\Phi$) along with the SM particles. Here the complex scalar field $\Phi$, which is singlet under SM but carries a non-zero $B - L$ charge, is required to break the $U(1)_{B-L}$ symmetry just above the inflation scale, and after the symmetry breaking the real part of $\Phi$ is identified as the inflaton in our scenario. Once this field $\Phi$ acquires the vacuum expectation value (vev), $U(1)_{B-L}$ symmetry is spontaneously broken around $10^{16}$ GeV, a scale which is governed by the scale of inflation as suggested by BICEP2 measurement [55]. The existence of the SM Higgs doublet ensures the breaking of electroweak symmetry at $\sim 246$ GeV.\(^2\) The important point to note here is that in such a scenario symmetry breaking of $U(1)_{B-L}$ and the SM electroweak vacuum take place at two very different energy scales.

\(^2\)We have defined the vev as $v/\sqrt{2}$. 
The part of the Lagrangian that contains the kinetic and potential terms of the scalars present in this theory is expressed as:

$$\mathcal{L} = (D_\mu S)^\dagger (D^\mu S) + (D_\mu \Phi)^\dagger (D^\mu \Phi) + m_s^2 (S^\dagger S) + m_\phi^2 (\Phi^\dagger \Phi) - \lambda_1 (S^\dagger S)^2 - \lambda_2 (\Phi^\dagger \Phi)^2 - \lambda_3 (S^\dagger S) (\Phi^\dagger \Phi),$$

(2.2)

where $S$ represents the SM Higgs field, $m_s$ and $m_\phi$ are the mass parameters of the Higgs and the U(1)$_{B-L}$ scalar fields respectively. Here, $\lambda_1$ and $\lambda_2$ are the quartic couplings of $S$ and $\Phi$ respectively and $\lambda_3$ is the coupling between $S$ and $\Phi$. In the unbroken SU(2)$_L \otimes U(1)_Y \otimes U(1)_{B-L}$ theory the covariant derivative is defined as:

$$D^\mu \equiv \partial^\mu + ig_2 \tau_j W_j^\mu + ig_1 Y B_1^\mu + ig Y^3 (g_T + g_{B-L} Q_{B-L}) B_2^\mu,$$

(2.3)

where $B_2^\mu$ is the $B - L$ charged gauge boson whose kinetic term is given as

$$\mathcal{L}^{B-L}_{K.E.} = \frac{1}{4} (\partial^\mu B_2^\nu - \partial^\nu B_2^\mu) (\partial_\mu B_{2\nu} - \partial_\nu B_{2\mu}).$$

(2.4)

Just above the inflationary scale, spontaneous breaking of the U(1)$_{B-L}$ symmetry yields $\Phi = v_0 + \phi(t, \mathbf{x})$ where $v_0 \equiv \sqrt{m_3^2 / \lambda_3}$ is the vev acquired by $\Phi$ and we have not written the phase part which yields the Goldstone mode. The real part, $\phi(t, \mathbf{x})$, of $\Phi$, apart from the vev, can be written as a background field $\phi_0(t)$ which plays the role of inflaton and fluctuations $\delta \phi(t, \mathbf{x})$ which give rise to the primordial perturbations during inflation. After the spontaneous breaking of $B - L$ symmetry, the scalar potential $V(S, \Phi)$ of this extended theory can be written as:

$$V(S, \Phi) = \lambda_1 (S^\dagger S)^2 - m_s^2 (S^\dagger S) + \lambda_2 \left( \Phi^\dagger \Phi - \frac{1}{2} v_0^2 \right)^2 + \lambda_3 (S^\dagger S) \left( \Phi^\dagger \Phi - \frac{1}{2} v_0^2 \right).$$

(2.5)

We see that various possible terms are generated in the scalar potential part of the Lagrangian, like $\lambda_3 v_0^2 S^\dagger S$, $\lambda_3 v_\phi S^\dagger S \phi_0$, $\lambda_3 S^\dagger S \phi_0 \phi_0$. The first term redefines the mass parameter of the $S$ field, the second term opens up the possibility of decay of inflaton into two SM Higgs fields during reheating. The third term introduces scattering of the light Higgs and the inflaton during the inflationary regime. We will concentrate on the importance of the second term later while discussing the decay of inflaton during reheating.

After the electroweak symmetry is broken at 246 GeV the Higgs field is redefined as $S = (0 \ 1/\sqrt{2} (v_s + s))^T$ and below this scale both the scalar fields get mixed and the physical fields ($\phi_l$ and $\phi_h$, where the subscripts $l$ and $h$ stands for ‘light’ and ‘heavy’ respectively) are achieved by diagonalizing the scalar mass matrix. The physical masses of these scalars are given as

$$M_{\phi_l, \phi_h}^2 = \frac{1}{2} \left[ \lambda_1 v_s^2 + \lambda_2 v_\phi^2 \mp \sqrt{(\lambda_1 v_s^2 - \lambda_2 v_\phi^2)^2 + \lambda_3^2 v_s^2 v_\phi^2} \right],$$

(2.6)

where the mixing angle is

$$\tan(2\alpha) = \frac{\lambda_3 v_s v_\phi}{\lambda_1 v_s^2 - \lambda_2 v_\phi^2}.$$  

(2.7)

\[\text{As our work includes both particle physics and cosmology where the Higgs field and the Hubble parameter, both conventionally represented by } H, \text{ have important roles to play, we choose the convention where the Hubble parameter is represented by } H \text{ and the SM Higgs field as } S \text{ to avoid confusion.}\]
We have set no Abelian mixing at tree level, i.e., \( g' = 0 \) at electroweak scale which can be done without any loss of generality. But this mixing will arise through the renormalisation group evolutions [56–58] and that has been taken into account in our analysis. Within this framework the gauge boson masses are

\[
M_Z^2 = \frac{1}{4} (g_1^2 + g_2^2) v_h^2, \quad M_{Z^0_{B-L}}^2 = 4 g_{B-L}^2 v_\phi^2.
\]  

(2.8)

One of our aims in this article is to seek for stability of electroweak vacuum all the way up to the Planck scale. We know the structure of the scalar potential is determined by the quartic couplings (\( \lambda_i \)) at large field values. Thus the \( \lambda_i \)s determine whether the potential is bounded from below or not. To achieve a stable vacuum the potential needs to be bounded from below at all energy scales, and this is a sufficient condition. Now, as stated before, the energy scales of the \( U(1)_{B-L} \) gauged extended SM and the low energy effective theory, which is just the SM, are well separated in our case. Thus the ‘decoupling theorem’ [59] states that the effects due to the new heavy particles of the extended theory would not affect the quartic coupling of the SM Higgs at low scales. This implies that at lower scales evolution of Higgs quartic coupling is governed by the SM particles only. Then it might indicate that extending the SM symmetry group would have no effect in stabilizing the electroweak vacuum at lower scales as the extended theory and the low energy effective SM theory are ‘decoupled’. But the advantage of extending the SM symmetry group with such a high scale \( U(1)_{B-L} \) gauge symmetry is that the extended theory contains a heavy scalar, which plays the role of an inflaton in our scenario. As has been proposed in [60], presence of a heavy scalar, besides the SM particles, eventually leads to a threshold correction to the SM Higgs quartic coupling and helps stabilize the electroweak vacuum as long as the mass of the heavy scalar lies below the instability scale of electroweak vacuum which is around \( 10^{10} \text{GeV} \). This is the key feature of our model and we would show that through one requires to break the \( U(1)_{B-L} \) symmetry at very high scale (\( \sim 10^{16} \text{GeV} \)) to have successful inflation at GUT scale, the mass of the inflaton can lie below the electroweak instability scale as the quartic coupling of the inflaton has to be fine tuned to yield the correct amplitude of scalar power spectrum, as we show below.

To show how the threshold correction, due to presence of a heavy scalar, modifies the evolution of Higgs quartic coupling \( \lambda_1 \) at lower scale [60], let us consider the scalar potential after \( U(1)_{B-L} \) symmetry breaking given in eq. (2.5). At lower energy scales, when the heavy scalar \( \Phi \) has reached its minima, its equation of motion yields

\[
\Phi^\dagger \Phi = \frac{1}{2} v_\phi^2 - \frac{\lambda_3}{2 \lambda_2} S^\dagger S.
\]  

(2.9)

Below the mass scale \( m_\phi \) of the inflaton, one can thus integrate out the heavy field \( \Phi \) using the above equation of motion and the potential given in eq. (2.5) becomes

\[
V(S)_{\text{eff}} = \left( \lambda_1 - \frac{\lambda_3^2}{4 \lambda_2} \right) (S^\dagger S)^2 - m_2^2 (S^\dagger S).
\]  

(2.10)

After that the dynamics of this theory is effectively governed by the SM particles where SM scalar potential is written as

\[
V(S)_{\text{SM}} \equiv \lambda_S (S^\dagger S)^2 - m_2^2 (S^\dagger S).
\]  

(2.11)
Here $\lambda_S$ is the SM Higgs quartic coupling related to the electroweak symmetry breaking scale and the SM Higgs mass only. At $m_\phi$ scale the impact of heavy inflaton field redefines the Higgs quartic coupling as $\lambda_S = (\lambda_1 - \frac{\lambda_2}{\Delta_m^2})$. This is a pure tree-level effect by which the heavy scalar of the extended theory affects the stability bound of the low energy effective theory even when these two theories are effectively decoupled. The Higgs quartic coupling $\lambda_S$ of the low energy effective theory receives a positive shift at the mass scale of the inflaton which thus helps avoid the instability which might have occurred above $m_\phi$ scale.

2.2 Inflation in the extended model

In this extended model under consideration the real part of the $U(1)_{B-L}$ breaking scalar field, i.e., $\phi_0$, plays the role of inflaton. Such a scenario has previously been considered in [48]. Before going into the details and particularities of our inflationary set up, we first briefly discuss what we know from the simplest single-field inflationary model in the light of recent BICEP2 as well as PLANCK observations. The amplitude of the two-point correlation function or the power spectrum of primordial scalar perturbations are measured through the two-point correlation of the temperature fluctuations in the CMBR. PLANCK has measured this value as [25]

$$P_R \sim 2.215 \times 10^{-9}. \quad (2.12)$$

The ratio of the tensor ($P_T$) and the scalar ($P_R$) power spectrum is represented as

$$r = \frac{P_T}{P_R}, \quad (2.13)$$

where $r$ is conventionally called the tensor-to-scalar ratio. This ratio $r$ has recently been measured by the BICEP2 experiment to be $0.20^{+0.07}_{-0.05}$ [3]. But after the release of PLANCK’s recent dust data [26] the observation of BICEP2 has been put under serious scrutiny. Though for the time being, before PLANCK and BICEP2 combine their observations, the upper-limit on $r$ set by PLANCK [26] still survives, i.e.,

$$r < 0.11 \quad (95\% \text{ CL}). \quad (2.14)$$

In single-field scenarios, the tensor power spectrum turns out to be a sole function of the Hubble parameter $H$ during inflation:

$$P_T = \frac{2}{\pi^2} \frac{H^2}{M_{Pl}^2}, \quad (2.15)$$

where $M_{Pl} \sim 2.4 \times 10^{18}$ GeV is the reduced Planck mass. As the Hubble parameter during inflation is related to the inflaton potential $V_\phi$,

$$H^2 = \frac{V_\phi}{3M_{Pl}^2}, \quad (2.16)$$

knowing $P_R$ and $r$ one can determine both the Hubble parameter during inflation

$$H \sim \sqrt{r} \times 10^{-4} M_{Pl}, \quad (2.17)$$
and the scale of inflation $V^{1/4}_\phi$

$$V^{1/4}_\phi \sim \left( \frac{r}{0.01} \right)^{1/4} \times 10^{16} \text{GeV}. \quad (2.18)$$

Furthermore, in a single-field model the scalar power spectrum turns out to be

$$P_R = \frac{H^2}{(2\pi)^2} \left( \frac{H^2}{\dot{\phi}_0^2} \right), \quad (2.19)$$

where the over-dot represents derivative with respect to cosmic time $t$ and this yields the tensor-to-scalar ratio as

$$r = \frac{8}{M^2_{Pl}} \left( \frac{\phi_0}{H} \right)^2 = \frac{8}{M^2_{Pl}} \left( \frac{d\phi_0}{dN} \right)^2, \quad (2.20)$$

where $N$ is the number of $e$-foldings during inflation. This indicates that the excursion of the inflaton field during inflation would be

$$\frac{\Delta \phi_0}{M_{Pl}} = \int_{N_{end}}^{N_{CMB}} dN \sqrt{\frac{r}{8}}, \quad (2.21)$$

where $N_{end}$ and $N_{CMB}$ are the number of $e$-foldings at the end of inflation and when the largest observable mode in the CMBR leaves the horizon before inflation ends, respectively. Assuming that $r$ would not change much during inflation, and $\Delta N \approx 65$ to solve the issues with Big Bang scenario, we have

$$\frac{\Delta \phi_0}{M_{Pl}} = \sqrt{530 \times r}. \quad (2.22)$$

Hence, for $r \geq O(10^{-2})$ the field excursion during inflation would be super-Planckian (large-field inflationary models), and for $r \leq O(10^{-2})$ it would be sub-Planckian (small-field inflationary models).

In the present model the inflaton potential can be written as [48]

$$V(\phi_0) = \frac{1}{4} \lambda_2 (\phi_0^2 - v_\phi^2)^2 + a \lambda_2 \log \left( \frac{\phi_0}{v_\phi} \right) \phi_0^4, \quad (2.23)$$

where we have

$$a \equiv \frac{1}{16\pi^2 \lambda_2} \left( 20 \lambda_2^2 + 2 \lambda_3^2 + 2 \lambda_2 \left( \sum_i (Y_{i_i}^{N_R})^2 - 24 g_{B-L}^2 \right) + 96 g_{B-L}^2 - \sum_i (Y_{i_i}^{N_R})^4 \right). \quad (2.24)$$

The above potential contains the radiative correction added to the tree-level one. Here $Y_{i_i}^{N_R}$ stand for the right handed neutrino Yukawa couplings. The value of ‘$a$’ determines whether the $U(1)_{B-L}$ symmetry is broken through the tree-level potential or the radiatively generated logarithmic term. As the value of ‘$a$’ mostly depends on the value of $g_{B-L}$ and $Y_{i_i}^{N_R}$, it can either be positive or negative depending upon the values of the couplings at inflationary scale. At tree level one can then identify the mass term of the inflaton as

$$m_\phi = \sqrt{\lambda_2 v_\phi}. \quad (2.25)$$
In large-field inflationary models one would naturally expect the quartic term with radiative corrections to dominate over the mass term in the inflaton potential and the form of the potential which would be responsible for inflation will be

$$V(\phi_0) \approx \frac{1}{4} \lambda_2 \phi_0^4 + a \lambda_2 \log \left( \frac{\phi_0}{\nu_0} \right) \phi_0^4.$$  \hfill (2.26)

The flatness of the potential is determined by the slow-roll parameters defined as

$$\epsilon_V = \frac{1}{2} M_{Pl}^2 \left( \frac{V'(V)}{V} \right)^2, \quad \eta_V = M_{Pl}^2 \left( \frac{V''(V)}{V} \right), \quad \xi_V^2 = M_{Pl}^4 \left( \frac{V'(V''(V))}{V^2} \right),$$  \hfill (2.27)

where the prime denotes derivative with respect to the field $\phi_0$. These slow-roll parameters remain small ($\epsilon_V, \eta_V \ll 1$) during inflation till $\epsilon_V$ becomes $\sim 1$, which marks the end of inflation. In a single-field scenario the tensor-to-scalar ratio $r$ and the scalar spectral index $n_s$ (which is a measure of the tilt of the scalar power spectrum) are related to the slow-roll parameters in the following way

$$r \approx 16 \epsilon_V, \quad n_s \approx 1 - 6 \epsilon_V + 2 \eta_V,$$  \hfill (2.28)

which are measured in CMBR observations. PLANCK measures the scalar spectral index as $n_s = 0.9603 \pm 0.0073$ [61]. The running of the scalar spectral index can also be determined in terms of the slow-roll parameters as

$$\frac{dn_s}{d \ln k} \approx 16 \epsilon_V \eta_V - 24 \epsilon_V^2 - 2 \xi_V^2.$$  \hfill (2.29)

If one just considers that quartic self-interaction of the inflaton potential, without the radiative term, to drive inflation, the the tensor-to-scalar ratio and the scalar spectral index turn out to be

$$r = \frac{128 M_{Pl}^2}{\phi_0^2}, \quad n_s = 1 - \frac{24 M_{Pl}^2}{\phi_0^2}.$$  \hfill (2.30)

The number of $e$-foldings can be calculated as

$$N_k = \frac{1}{M_{Pl}^2} \int_{\phi_0_{end}}^{\phi_k} \frac{V(\phi)}{V'} d\phi,$$  \hfill (2.31)

where $\phi_{0_k}$ is the field value at the co-moving scale $k$ and $\phi_{end}$ is the field value at the end of inflation. This yields

$$N_k \approx \frac{\phi_{0_k}^2}{8 M_{Pl}^2}.$$  \hfill (2.32)

where $k_s$ denotes the pivot scale and it is considered that $\phi_{0_{end}} \ll \phi_{0_s}$. Hence, if the mode corresponding to the pivot scale would have left the horizon around 65 $e$-foldings before inflation ends, then the above expression helps to determine the field value during that time,
which turns out to be $\phi_0, \sim 23 M_{\text{Pl}}$. The tensor-to-scalar ratio and the scalar spectral index at scale $k$ can be expressed in terms of the $e$-foldings as

$$r_k = \frac{16}{N_k}, \quad n_{s_k} = 1 - \frac{3}{N_k}.$$  \hspace{1cm} (2.33)

The number of $e$-folds when the pivot scale crosses the horizon during inflation can also be expressed as

$$N_e \simeq 65 + 2 \ln \left( \frac{V_\phi(\phi_0)}{10^{14} \text{ GeV}} \right) - \ln \left( \frac{T_f}{10^{10} \text{ GeV}} \right),$$  \hspace{1cm} (2.34)

where $T_f$ is the temperature at the end of inflation which one can consider as the reheat temperature $T_{\text{RH}}$. If the pivot scale set by PLANCK, i.e., $k_\ast = 0.002 \text{ Mpc}^{-1}$, crosses the horizon during inflation when $N_e \sim 65$ then it yields large tensor-to-scalar ratio as $r_\ast \sim 0.25$ which is also large enough even for BICEP2 observations. This corresponds to the field excursion during inflation to be $\Delta \phi \sim 12 M_{\text{Pl}}$. Hence our aim now would be to generate lower values of $r$ while keeping the scenario consistent with the observations of $n_s$ and $\mathcal{P}_R$ by PLANCK.

It has been pointed out in [62] that the radiative corrections to the quartic potential play an important role in lowering the tensor-to-scalar ratio. Hence, for our inflationary scenario we consider the inflaton potential including radiative correction for inflation. When inflation is driven by this quartic potential, we find

$$V_\phi = \frac{1}{4} \lambda_2 \phi_0^4 \left[ 1 + 4a \ln \left( \frac{\phi_0}{v_\phi} \right) \right],$$

$$V'_\phi = \lambda_2 \phi_0^3 \left[ 1 + a + 4a \ln \left( \frac{\phi_0}{v_\phi} \right) \right],$$

$$V''_\phi = 3 \lambda_2 \phi_0^2 \left[ 1 + \frac{7}{3}a + 4a \ln \left( \frac{\phi_0}{v_\phi} \right) \right],$$

$$V'''_\phi = 6 \lambda_2 \phi_0 \left[ 1 + \frac{13}{3}a + 4a \ln \left( \frac{\phi_0}{v_\phi} \right) \right].$$

These give the slow-roll parameters as:

$$\epsilon_V = \frac{8 M_{\text{Pl}}^2}{\phi_0^2} \left[ \frac{u^2}{(u-1)^2} \right], \quad \eta_V = \frac{12 M_{\text{Pl}}^2}{\phi_0^2} \left[ \frac{u + 4/3}{u-1} \right], \quad \xi_V = \frac{96 M_{\text{Pl}}^4}{\phi_0^4} \left[ \frac{u(u + 10/3)}{(u-1)^2} \right].$$

where we have defined $u = (1 + a + 4a \ln \phi_0/v_\phi)/a$. Hence the tensor-to-scalar ratio, the scalar spectral index and the running of the scalar spectral index can be written as:

$$r = \frac{128 M_{\text{Pl}}^2}{\phi_0^2} \frac{u^2}{(u-1)^2},$$

$$n_s = 1 - \frac{8 M_{\text{Pl}}^2}{\phi_0^2} \frac{3u^2 - u + 4}{(u-1)^2},$$

$$\frac{dn_s}{d \ln k} = - \frac{64 M_{\text{Pl}}^4}{\phi_0^4} \left[ \frac{u(3u^3 - 4u^2 + 15u + 10)}{(u-1)^4} \right].$$
respectively. If we consider the radiative correction in the scalar potential given in eq. (2.35) is negligible which implies $a \to 0$, i.e., $u \to \infty$, the standard expressions for $r$ and $n_s$ for quartic coupling can be recovered. Furthermore, the scalar power spectrum, given in eq. (2.19), can also be written as

$$P_R = \frac{1}{24 \pi^2 \epsilon_V} \left( \frac{V_\phi}{M_{Pl}^4} \right) = \frac{1}{12 \pi^2 M_{Pl}^2} \frac{V_\phi^3}{V''_\phi},$$  \hspace{1cm} (2.38)$$

where in the first equality we have used the Hubble slow-roll parameter $\epsilon = \frac{1}{2 M_{Pl}^2} \frac{\dot{\phi}^2}{H^2}$ and the Friedman equation during inflation given in eq. (2.16) with $\epsilon \approx \epsilon_V$. The power spectrum with radiatively corrected potential term turns out to be

$$P_R = \frac{\lambda_2^2}{768 \pi^2} \left( \frac{\phi_0}{M_{Pl}} \right)^6 \frac{a(u-1)^3}{u^2}. \hspace{1cm} (2.39)$$

Now, this scenario can be realised in two cases. In the limit when $u \gg 1$, one can have $|a| \ll 1$, when the radiative corrections become negligible. In such a case the standard results for $\phi^4$ potential should be retrieved. The other branch is when $u \approx 1$ which gives $a \sim -(4 \ln(\phi_0/v))^{-1}$. This branch is called the hill-top solution.

We also require to determine the reheat temperature in order to compute the number of e-foldings which corresponds to the pivot scale as given in eq. (2.34). We notice that, apart from the self-interaction term, the inflaton field is also coupled to the SM Higgs field via the mixing term $\lambda_3$ which allows it to decay into a pair of SM Higgs during inflation. The decay rate of such an interaction is given as [48]:

$$\Gamma_S(\phi_0 \to SS) = \frac{\lambda_3^2 v_\phi^2}{32 \pi m_{\phi^2}}. \hspace{1cm} (2.40)$$

This decay of inflaton field into SM Higgs would make inflaton unstable for larger values of $\lambda_3$. Thus one requires to restrain the decay width of the inflaton during inflation. This requirement can be met if one demands that $\Gamma_S < m_\phi$ which yields

$$\lambda_3 < \sqrt{32 \pi \lambda_2}. \hspace{1cm} (2.41)$$

From eq. (2.40) we can also roughly estimate the order of reheating temperature $T_{RH}$ if the reheating phase is dominated by the Higgs decay. If during the reheating phase the inflaton and its decay products are just in equilibrium then $\Gamma_S \sim H$ where $H$ is the Hubble parameter during the radiation dominated reheating phase. This condition yields

$$\frac{\lambda_3^2 v_\phi^2}{32 \pi m_{\phi^2}} = \frac{\sqrt{\pi^2}}{90} \frac{g_*}{M_{Pl}^2} T_{RH}^2, \hspace{1cm} (2.42)$$

where $g_* \sim 100$.

Now, let us determine the parameters for a large-field inflationary scenario and take $\phi_{0,s} \sim 23 M_{Pl}$. Putting the central value of scalar spectral index as $n_s = 0.9603$ one gets two solutions for $u$ at the pivot scale as $u_s = -0.333$ and $u_s = -11.001$. The first solution indicates a hilltop branch inflation whereas the second solution gives rise to a $\phi^4$-branch inflation. Let us now analyze the parameters for these two scenarios:
\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{potential.png}
\caption{Forms of possible inflaton potentials in the present scenario.}
\end{figure}

- **Hilltop inflation:** if one sets the vev of the inflaton as $10^{16}$ GeV, then for $u_s = -0.333$ one gets $a_s \sim -0.028$. This indicates that the field value at the end of inflation, when $\epsilon_V \approx 1$, would be $\phi_{0_{\text{end}}} \sim 0.71 M_{\text{Pl}}$. This value of $u_s$ yields the tensor-to-scalar ratio as $r_s = 0.015$ and the inflaton quartic coupling, from the observation of the scalar power amplitude by PLANCK, as $\lambda_2 \sim 1.89 \times 10^{-13}$. This then yields the tree-level mass of the inflaton as $m_\phi \sim 4.3 \times 10^9$ GeV. The running of the spectral index in such a scenario would be $\frac{d n_s}{d \ln k} \sim 1.07 \times 10^{-4}$. In this scenario the inflaton-Higgs coupling can be of the order of $\sim 10^{-6}$, which yields the reheat temperature as $T_{RH} \sim 1.29 \times 10^{13}$ GeV. This reheat temperature and the energy-scale of inflation yield the e-folding at which pivot scale would have exited the horizon as $N_s \sim 67$. Figure 1(a) shows the form of the potential for the hilltop inflation.

- **$\phi^4$-branch inflation:** if one sets the vev of the inflaton as $10^{16}$ GeV, then for $u_s = -11.001$ one gets $a_s \sim -0.022$, which is negligible. This indicates that the field value at the end of inflation, when $\epsilon_V \approx 1$, would be $\phi_{0_{\text{end}}} \sim 2.6 M_{\text{Pl}}$. This value of $u_s$ yields the tensor-to-scalar ratio as $r_s = 0.203$ and the inflaton quartic coupling, from the observation of the scalar power amplitude, as $\lambda_2 \sim 3.6 \times 10^{-13}$. This then yields the tree-level mass of the inflaton as $m_\phi \sim 6.0 \times 10^9$ GeV. The running of the spectral index in such a scenario would be $\frac{d n_s}{d \ln k} \sim -5.6 \times 10^{-4}$. In this scenario the inflaton-Higgs coupling can be of the order of $\sim 10^{-6}$, which yields the reheat temperature as $T_{RH} \sim 1.09 \times 10^{13}$ GeV. This reheat temperature and the energy-scale of inflation yield the e-folding at which pivot scale would have exited the horizon as $N_s \sim 67$. Figure 1(b) shows the form of the potential for the $\phi^4$-branch inflation.

### 3 Results

In the previous section we have described the particle content of the U(1)$_{B-L}$ extended SM, the dynamics of its scalar sector and the framework of how inflation can be included in this model. Our aim is to stabilize the vacuum all the way up to the Planck scale and also to accommodate inflationary paradigm successfully. In this section we would show how the aspect of vacuum stability can be achieved in our framework if inflation is driven by the inflaton’s radiatively corrected quartic self-interaction term.
3.1 Vacuum stability

In the previous subsection, we have shown that to achieve successful inflation, both the inflaton quartic coupling and the interaction quartic coupling have to be fine tuned. Finetuning of inflaton quartic coupling evidently brings down the mass scale of the inflaton field which turns out to be below the instability scale of the electroweak vacuum. Following [60], one can integrate out the heavy inflaton field below its mass scale which then adds a tree-level threshold correction to the low energy effective Higgs quartic coupling $\lambda_S$ as (see eq. (2.10))

$$\lambda_1 = \lambda_S + \frac{\lambda_3^2}{4\lambda_2}. \quad (3.1)$$

Hence below the inflaton mass scale the stability condition ($\lambda_S > 0$) for the SM Higgs quartic coupling would get shifted upwards $\lambda_1 > \delta \lambda \equiv \frac{\lambda_3^2}{4\lambda_2}$. The other two quartic couplings $\lambda_2$ and $\lambda_3$ would start evolving at energies above this mass scale. The relevant RGEs are written at the appendix A.1. To illustrate the threshold effect, the running of Higgs quartic coupling has been shown in figure 2 for both the inflationary scenarios. Here we have taken the heavy scalar mass $m_\phi = 4 \times 10^9$ GeV before which $\lambda_S$ runs according to SM $\beta$-functions. Beyond that point new loop effects due to the extended theory start to affect its running and the discrete jump in the Higgs quartic coupling at $m_\phi$ is due to the threshold effect.

Apart from the SM fermions this model also contains three right handed neutrinos, $N_{Ri}$, which appear in the Lagrangian as

$$-\mathcal{L}_Y = Y_{\nu i} \bar{l}_i \tilde{S} N_{Rj} + Y_{ij}^{N_R} \bar{(N_R)^c} N_{Rj} \Phi + h.c., \quad (3.2)$$

where $\tilde{S} = i\sigma_2 S^*$, $\sigma_2$ being the second Pauli matrix and $l$ stands for the SM lepton doublet. The second term in the above Lagrangian gives rise to the coupling of the inflaton to heavy
right handed neutrinos and also masses for $N_R$. It is important to note that when the $(B - L)$ symmetry is broken at the TeV scale the masses of the right handed neutrinos are less compared to the present scenario. In case of TeV scale breaking the Yukawa couplings ($Y^{\nu L}$) giving rise to the Dirac mass of light neutrinos have to be vanishingly small unless some special textures are considered. Thus in such cases, impact of $Y^{\nu L}$ in the evolutions of the quartic and other necessary couplings is negligible. But in the present case the right handed neutrino masses are very heavy $\sim 10^{11-13}$ GeV, due to high $U(1)_{B-L}$ breaking scale. Thus the light neutrino masses are still light $\sim \mathcal{O}$(eV) even with $Y^{\nu L} \sim \mathcal{O}(1)$. Hence unlike the cases, where $U(1)_{B-L}$ symmetry is broken at TeV scale, one can not ignore the contributions of light neutrino Yukawa couplings $Y^{\nu L}$ in the RGEs in our scenario (see appendix A.2).

### 3.2 Other constraints

Looking at the threshold correction, given in eq. (3.1), which is essential for electroweak vacuum stability, it may seem that $\lambda_3 < 0$ can still be retained as a possible condition. But, in our analysis this opportunity of achieving larger parameter space for $\lambda_3$ is restricted as here $\lambda_2$ is very small $\sim 10^{-13}$ due to inflationary constraints. The absolute value of $\lambda_3$ can never be too large as it affects the running of $\lambda_2$ by driving its value to a much larger value which might not be able to explain inflationary dynamics. Thus $\lambda_3$ is constrained from above by the requirement of inflation. The smallness of $|\lambda_3|$ ensures that the two scalars present in the theory are basically decoupled from each other as the mixing angle then becomes too small, see eq. (2.7). This confirms that the ‘decoupling theorem’ holds good in our scenario.

At the inflaton mass scale from which running of $\lambda_2$ and $\lambda_3$ become important, explicit values of $\lambda_3$ and $\lambda_2$ have been set as $8.0 \times 10^{-8}$ and $5 \times 10^{-13}$ ($2.7 \times 10^{-13}$) for hilltop inflation ($\phi^4$-branch inflation) respectively. These values ensures both vacuum stability up to the Planck scale and inflation with correct scalar amplitude at GUT scale. The $B - L$ gauge coupling, $g_{B-L}$, also affects the running of $\lambda_2$ from $m_{\phi}$ scale to inflation scale (see appendix (A)). The radiative correction to the scalar potential is shown in eq. (2.24) which depends on both $\lambda_3$ and $g_{B-L}$. We set $g_{B-L}$ at $\sim 10^{-5}$ and show the variation of $a$ with energy scale in figure 3 for both the inflationary scenarios.

In a previous discussion we have presented a possible mode in which the inflaton can decay into a pair of Higgs fields during reheating. In passing, we note another possible way of reheating in our model. The inflaton can also decay to two heavy right handed neutrinos where the decay rate is given by

$$\Gamma_N(\phi_0 \to N_{Ri}N_{Rj}) = \frac{(Y^{NR}_{i})^2}{64\pi} m_\phi.$$  \hspace{1cm} (3.3)

But this $\Gamma_N$ is much smaller than $\Gamma_S(\phi_0 \to SS)$ (as $Y^{NR}_{i} \sim 10^{-4}$ in our scenario) yielding a much smaller reheating temperature compared to $T_{RH}$ when the inflaton prominently decays into Higgs. In general if both the decay channels are open for the inflaton, reheating will prominently happen through the Higgs decay of the inflaton.

### 3.3 Quadratic divergence of Higgs mass and the Veltman criteria

We would like to note in passing that the stabilization of the SM Higgs boson mass under the quadratic divergences specially in high scale theory is an unavoidable issue. The generic problem with any high scale non-supersymmetric models is related to the stabilization of scalar masses, specially the SM Higgs mass. Due to the quadratic divergences the SM Higgs
mass acquires a correction proportional to $\Lambda^2$, where $\Lambda$ is the scale of new physics. In supersymmetric theory these corrections automatically get cancelled out with the ones coming from their supersymmetric partners. To avoid such large contributions to scalar mass in non-supersymmetric models one needs to impose the Veltman condition. This prescription, as suggested in [63, 64], confirms the removal of quadratic divergences of the scalar masses to stabilize them. We note here that in our model the scalar masses might obtain a large radiative correction. To avoid such catastrophe in scalar masses we need to satisfy the Veltman criteria (VC), which for this scenario would be [65]

$$\delta m_s^2 \propto v_\phi^2[(2\lambda_1 + \lambda_3/3) \cos^4 \alpha + \left(\frac{2M_W^2 + M_Z^2}{v^2}\right) \cos^2 \alpha$$

$$+ 4g_{B-L}^2 \sin^2 \alpha - 4(Y_t^2 \cos^2 \alpha + (Y_{\nu L}^\nu)^2 \cos^2 \alpha + (Y_{N R}^N)^2 \sin^2 \alpha)], \quad (3.4)$$

with $\cos^2 \alpha \sim 0.99879$, where $\alpha$ is the mixing angle as given in eq. (2.7). Since $v_\phi$ is very large for our case, $Z_{B-L}$ and $N_R$ will not affect this criteria much. Also $\lambda_3$ is very small. But the light neutrino Yukawa coupling can be large here and thus its impact can be sizeable. This has been added with top quark contribution. We find that within the available parameter space in our model it is indeed possible to satisfy VC either at the inflation scale or at Planck scale. With light neutrino Yukawa to be 0.1462 and 0.2413 the VC can be satisfied at the inflation and the Planck scales respectively. This implies that the light Higgs remains light at the inflation scale and the decay of the inflaton, considered in this paper for explaining the reheating, does not suffer any catastrophe under the impact of radiative corrections. Thus we can stabilize the Higgs mass at the inflation scale but perhaps this mechanism does not solve the stabilization at other scales.

Figure 3. Variation of $a$, see eq. (2.24), with respect to the energy scale.
4 Discussion and conclusion

The recent discovery of SM Higgs scalar and the detection of B-mode polarization of CMBR by BICEP2, if stands the test of time, are quite remarkable as these observations would help understand the evolution of our universe. Though these two observations carry signatures of physics taking place at very different scales, we tried to connect them within a single particle physics picture in this article. The observation by BICEP2 of the CMBR B-mode polarization, if confirmed, would also indicate the scale of inflation to be around $\sim 10^{16}$ GeV. This energy scale is quite interesting in the field of particle physics as incidentally this scale turns out to be the unification scale. In this work, we have adopted a gauge extended SM scenario which contains a SM singlet scalar field with a new $U(1)_{B-L}$ gauge charge. This field acquires vev at a very high scale and breaks the $U(1)_{B-L}$ symmetry spontaneously and also couples to the SM particles. Apart from this SM singlet scalar, there are three right handed neutrinos which successfully generates the light neutrino masses through type-I seesaw without fine-tuning the Dirac Yukawa coupling.

The electroweak breaking scale is around 246 GeV whereas in this case the $U(1)_{B-L}$ breaking scale should lie near the GUT scale so that such high-scale inflation can take place as has been demanded by BICEP2. But, as the $U(1)_{B-L}$ and electroweak breaking scales lie far apart in our scenario, these two theories are basically decoupled from each other, as has been demanded by the ‘decoupling theorem’. Then it might imply that extending the SM by such high scale $U(1)_{B-L}$ gauge theory fails to serve its purpose of taking care of the stability of electroweak vacuum. But the advantage of introducing such high scale $U(1)_{B-L}$ symmetry is that it provides a heavy scalar $\Phi$ in the theory, whose real part plays the role of the inflaton in such a scenario. Presence of a heavy scalar yields a threshold correction to the Higgs quartic coupling, if one integrates out this heavy scalar below its mass scale. Hence if the mass of this heavy scalar lies below the electroweak instability scale ($\sim 10^{10}$ GeV), the threshold correction eventually helps avoid the instability of the vacuum by correctly uplifting the value of the SM Higgs quartic coupling at this mass scale. Hence the key point is to keep the mass scale of the heavy scalar in the theory below the electroweak scale if one wants the threshold corrections to help stabilize the vacuum, even though the explicit value of the mass of the heavy scalar does not play any important role.

Even at this point, it might seem that the heavy scalar mass, which is basically the inflaton mass in our case, would lie above the electroweak instability scale as the $U(1)_{B-L}$ breaking scale lies far above this instability scale. Thus it might seem that even the threshold effect would not keep the vacuum stable in such a case. The important point to note here is that the heavy scalar, which yields the threshold correction to the SM Higgs quartic coupling, also plays the role of the inflaton in such a case. In the large field inflation scenario, supported by BICEP2 observations, inflation is naturally driven by the radiatively corrected quartic potential of the inflaton. Thus the inflaton’s quartic coupling parameter, which eventually determines its mass, has to be extremely fine-tuned so that the amplitude of the scalar power spectrum remains in accordance with observation. Such high fine-tuning of inflaton quartic coupling evidently brings down the mass of the inflaton below the electroweak vacuum instability scale, enabling the threshold effect to keep the electroweak vacuum stable.

The radiatively corrected inflaton potential can also produce much lower tensor-to-scalar ratio $r$ in the hill-top branch inflation scenario, which remains much below the upper-limit set by PLANCK and hence observationally viable. The other scenario, namely $\phi^4$-branch inflation, produces $r \sim 0.203$ which is compatible with BICEP2 observations. Both these
scenarios produce very low running of the scalar spectral index ($\mathcal{O}(10^{-4})$). Thus the $\phi^4$-branch inflation scenario would not resolve the discrepancy between PLANCK and BICEP2 observation of $r$, while the hilltop inflationary scenario is observationally safe if the BICEP2 observations are ruled out in future.

With such a set-up, we have computed the correlated bounds on scalar quartic couplings and Yukawa couplings compatible with vacuum stability and perturbativity all the way up to the Planck scale. In order to have successful inflation in this particle physics model, we showed that fine-tuning of $U(1)_{B-L}$ quartic couplings allow a chaotic-type inflation driven by the radiatively corrected quartic potential of the real part of the SM singlet scalar and the scenario remains in accordance with the observations of BICEP2 and PLANCK. We are also able to ascertain a particular parameter space where the criteria for vacuum stability and requirements to explain inflationary dynamics are simultaneously satisfied. To our advantage, this model also yields reheating after inflation through the channel where the inflaton possibly decays into a pair of SM Higgs field. Constraints coming from such reheating scenario are automatically satisfied within our choice of parameter space. Hence to summarize, the $U(1)_{B-L}$ symmetry extended SM considered here is capable of

- stabilizing SM electro-weak vacuum all the way up to the Planck scale,
- yielding light neutrino masses without fine-tuning of Dirac Yukawa couplings,
- accommodate large-field chaotic-type inflation which is in accordance with BICEP2 and PLANCK’s observations,
- successfully reheat the universe at the end of inflation.

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A Renormalisation of couplings

Here we have depicted the necessary renormalisation group (RG) evolutions for $U(1)_{B-L}$ extended Standard Model as suggested in [21, 66]. The necessary parameters are: scalar quartic ($\lambda_{1,2,3}$), top Yukawa ($Y_t$), Dirac Yukawa ($Y^{\nu L}$), right handed neutrino Yukawa ($Y^{\nu R}$), gauge ($g_{1,2,3}, g_{B-L}, g'$) couplings.

A.1 RG equations for scalar quartic couplings

Here we have illustrated the RGEs for $\lambda_{1,2,3}$. Their evolutions with scales involve the renormalisation of other parameters, see eqs. (A.2), and A.3.

\[
(16\pi^2) \frac{d\lambda_1}{dt} = \left[ 24\lambda_1^2 + \lambda_2^2 - 6Y_t^4 + \frac{9}{8}g_2^4 + \frac{3}{8}g_1^4 + \frac{3}{4}g_1^2g_2^2 + \frac{3}{4}g_2^4 + \frac{3}{4}g_1^2g' + \frac{3}{4}g_2^2g' + 18\lambda_1g_2^2 + 12\lambda_1Y_t^2 - 3\lambda_1g_2^2 - 3\lambda_1g_1^2 - 2(Y^{\nu L})^4 + 4\lambda_1(Y^{\nu L})^2 \right],
\]

\[
(16\pi^2) \frac{d\lambda_2}{dt} = \left[ 20\lambda_2^2 + 2\lambda_3^2 - \text{Tr}[(Y^{\nu R})^4] + 96g_{B-L}^4 + 8\lambda_2^2 \text{Tr}[(Y^{\nu R})^2] \right],
\]
\[(16\pi^2) \frac{d\lambda_3}{dt} = \lambda_3 \left[ 12\lambda_1 + 8\lambda_2 + 4\lambda_3 + 6Y_t^2 - \frac{3}{2}(3g_2^2 - g_1^2 - g'^2) + 4\text{Tr}[(Y_N^R)^2] - 24g_{B-L}^2 + 12g_{B-L}^2g' \right]. \quad (A.1)\]

### A.2 RG equations for Yukawa couplings

Here we have characterized the RGEs of different Yukawa couplings that affect the evolutions of quartic couplings as given in eq. (A.1).

\[(16\pi^2) \frac{dY_t}{dt} = Y_t \left[ \frac{9}{2}Y_t^2 - 8g_3^2 - \frac{9}{4}g_2^2 - \frac{17}{12}(g_1^2 + g_2^2 + g'^2) - \frac{2}{3}g_{B-L}^2 - \frac{5}{3}g_{B-L}g' + (Y_{\nu L})^2 \right], \quad (A.2)\]

\[(16\pi^2) \frac{dY^N_R}{dt} = Y^N_R \left[ 4(Y^N_R)^2 + 2\text{Tr}[(Y^N_R)^2] - 6g_{B-L}^2 \right], \quad (A.2)\]

\[(16\pi^2) \frac{dY_{\nu L}}{dt} = Y_{\nu L} \left[ \frac{5}{2}(Y_{\nu L})^2 + 3Y_t^2 - \frac{9}{4}g_2^2 - \frac{3}{4}g_1^2 - 6g_{B-L}^2 \right]. \quad (A.2)\]

### A.3 RG equations for the gauge couplings

The gauge coupling runnings are delineated here as they also play crucial role while adjudging the RGEs of quartic couplings and Yukawa couplings, see eqs. (A.1), and A.2.

\[(16\pi^2) \frac{dg_3}{dt} = \left[ -\frac{7}{6} \right] g_3^3, \]
\[(16\pi^2) \frac{dg_2}{dt} = \left[ -\frac{19}{2} \right] g_2^3, \]
\[(16\pi^2) \frac{dg_1}{dt} = \left[ \frac{41}{6} \right] g_1^3, \quad (A.3)\]
\[(16\pi^2) \frac{dg_{B-L}}{dt} = \left[ 12g_{B-L}^3 + (32/3)g_{B-L}g' + (41/6)g_{B-L}g'^2 \right], \]
\[(16\pi^2) \frac{dg'}{dt} = \left[ \frac{41}{6} (g'^3 + 2g_1^2g') + \frac{32}{3}g_{B-L}(g'^2 + g_1^2) + 12g_{B-L}g' \right]. \]

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