A minimal gauge inflation model *

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Abstract: In this paper, we present a gauge inflation model based on the orbifold $M_4\times S^1/\mathbb{Z}_2$ with non-Abelian SU(2) gauge symmetry, which is probably the simplest model in this category. As the inflaton potential is fully radiatively generated exclusively by gauge self-interactions, the model is predictive; thus, it is protected by gauge symmetry itself, without the introduction of any additional matter fields or arbitrary interactions. We show that the model fully agrees with the recent cosmological observations within the controlled perturbative regime of gauge interactions, $g_4\lesssim 1/(2\pi RM_{\rm P})$, with the compactification radius ($10\lesssim RM_{\rm P}\lesssim 100$): the expected magnitude of the curvature perturbation power spectrum and the value of the corresponding spectral index are in perfect agreement with the recent observations. The model also predicts a large fraction of the gravitational waves, negligible non-Gaussianity, and a sufficiently high reheating temperature.

Keywords: inflation, non-Gaussianity, extra dimension

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1 Introduction

It is now widely accepted that the theory of the early accelerating expansion of the Universe or cosmological inflation [1], has the capacity to solve several cosmological problems including the problem of flatness, horizon, and magnetic monopole and can also provide the required initial conditions for the subsequent hot Big Bang evolution of the observed Universe [2]. From a particle physics point of view, inflation takes place due to one or multiple number of scalar fields, known as inflatons, which dominate the energy density of the Universe during the inflationary period, with a flat potential [3]. Under such a condition, a nearly scale invariant curvature perturbation $\mathcal{R}$ is produced, according to the precise measurements of the anisotropies of the cosmic microwave background (CMB) and observations of the large-scale structure [4]. The most important recent observation on the anisotropy of CMB is from the Planck observatory [5]. It should be noted, that the observation of the primordial gravitational waves from the inflationary period, which can be expected soon, provides a valuable testing ground for the theoretical realization of inflation (see, e.g. the Higgs inflation [6]).

For the theoretical realization of inflation, symmetry principles are often introduced to ensure that the inflaton potential is flat. The shift symmetry, under which the Lagrangian is invariant with respect to the translation of the inflaton field by a constant amount as $\phi\to\phi+a$, is one of such symmetries. As long as the shift symmetry is unbroken, the potential is completely flat; however, when the symmetry is broken at a scale f, the pseudo-Nambu-Goldstone boson(pNGB) acquires a potential with a certain tilt, which is controlled by the scale f as in the natural inflation model, where the inflaton potential is given as [7]:

$$V(\phi)=\Lambda^4\left[1\pm\cos\left(\frac{\phi}{f}\right)\right]. \tag{1}$$

To successfully apply this potential to inflation, a super-Planckian scale, $f>M_{\rm P}$, is required; however, this large scale itself has been regarded as an essential drawback of the model (see Ref. [8]).

The framework of higher-dimensional gauge theory can solve the super-Planckian problem [9]. The inflaton dynamics is controlled by the gauge-invariant Wilson line of the higher-dimensional gauge theory:

$$e^{i\theta}\equiv\exp\left(ig\oint A_5 dy\right), \tag{2}$$

where g is the gauge coupling constant and A_5 is the extra dimensional component of the higher-dimensional

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gauge field. In the existing models, higher-dimensional matter fields are additionally introduced to generate the proper form of the inflaton potential [9, 10]. Applying these matter fields at one-loop level, the generated potential has essentially the same form as that in in Eq. (1), while in this case, the scale f is determined by the gauge coupling constant and the compactification radius $f=1/(g_4R)$. It is important to note that the potential can be trusted, because f is a derived scale, effectively parametrizing the inflationary dynamics *within* the perturbative regime of the coupling constant: $g_4 \lesssim 1/(RM_{\rm P}) \ll \sqrt{4\pi}$.

In this paper, we exploit the benefits of this line study and suggest a minimal realistic model, which enables definite predictions for cosmological observations without assuming ad hoc matter fields and arbitrarily chosen interactions between them. The model is based on a gauge theory in the simplest orbifold extra dimension $S^1/\mathbb{Z}_2$, and a non-Abelian gauge group SU(2) of rank 1. The gauge self-interaction, without introducing any matter fields, induces the inflaton potential, which is completely determined by only two parameters: the size of the extra dimension, R, and the gauge coupling constant, g.

The rest of the paper is organized as follows: In the next section, we provide a detailed description of the theoretical setup. Section 3 is devoted to the cosmological predictions of the model; in particular, we provide the analytic expressions for the essential observable quantities. Finally, we address the reheating process, followed by the conclusion in Section 4.

2 The model

As mentioned in the previous section, the introduction of any exotic matter must be avoided only for inflation. The gauge self-interactions of non-Abelian gauge theory can generate the potential by the quantum effects; thus, we applied an SU(2) symmetry of rank 1 without introducing any other exotic matters. The extra dimension is assumed to be compactified by an $S^1/\mathbb{Z}_2$ orbifold [1)].

The SU(2) gauge theory on the orbifold is constructed by specifying two independent parity conditions at the two fixed points, $y=0$ and $y=\pi R$, where R is the compactification radius, as

$$A_\mu(x,-y)=P_0A_\mu(x,y)P_0, \tag{3}$$

$$A_5(x,-y)=-P_0A_5(x,y)P_0, \tag{4}$$

$$A_\mu(x,\pi R-y)=P_1A_\mu(x,\pi R+y)P_1, \tag{5}$$

$$A_5(x,\pi R-y)=-P_1A_5(x,\pi R+y)P_1, \tag{6}$$

where P_0 and P_1 are 2×2 matrices satisfying $P_0^2=P_1^2=1$. The translational transformation, $y\to y+2\pi R$, is generated by successive operations of the parity operators, P_1P_0. Assuming that $P_0=P_1=\mathrm{diag}(1,-1)$, the SU(2) gauge symmetry is reduced to U(1) by the orbifold projection at the classical level. Here, we explicitly write out the parity assignment with P_0 and P_1 as

$$A_\mu=\begin{pmatrix}(++) & (--)\\ (--) & (++)\end{pmatrix}, \tag{7}$$

$$A_5=\begin{pmatrix}(--) & (++)\\ (++) & (--)\end{pmatrix}; \tag{8}$$

thus, the zero modes are given by the (+,+) boundary conditions, hence A_μ^3 and $A_5^{1,2}$ are the zero modes. The component field A_M^a belongs to the $\sigma_a/2$ of the SU(2) generator. The scalar field $A_5^{1,2}$ can develop a vacuum expectation value, which can be written in the form $A_5^a\sim(\phi,0,0)$, due to the remaining U(1) global symmetry. Considering the effects of the gauge, ghost, and scalar-self interactions, the one-loop effective potential for the field ϕ can be calculated as

$$V_{\text{1-loop}}(\phi)=-\frac{9}{(2\pi)^6R^4}\sum_{n=1}^{\infty}\frac{\cos(n\phi/f_{\rm eff})}{n^5}, \tag{9}$$

where the effective decay constant is introduced as

$$f_{\rm eff}\equiv\frac{1}{\sqrt{2\pi R}g}=\frac{1}{2\pi g_4R} \tag{10}$$

for the canonical normalization of ϕ [11] (see also [13]). Here, we can add a cosmological constant

$$\frac{9\zeta(5)}{(2\pi)^6R^4}, \tag{11}$$

where $\zeta(5)=\sum_{n=1}^{\infty}n^{-5}$, such that $V_{\text{1-loop}}(0)=0$ to fix the cosmological constant. Finally, the total inflaton potential is given by the sum of Eqs. (9) and (11) as:

$$V(\phi)=\frac{9}{(2\pi)^6R^4}\sum_{n=1}^{\infty}\frac{1}{n^5}\left[1-\cos\left(\frac{n\phi}{f_{\rm eff}}\right)\right]. \tag{12}$$

In principle, we can introduce additional matter fields; however, for simplicity and predictability we refrain to do so. In the next section, we calculate the observable quantities based on the above potential and provide analytic expressions, which can be useful in the future analysis of the similar models.

3 Cosmological evolution

In this section, we study the cosmological evolution of the model described in the previous section in detail. For analytic simplicity, we consider only the first term of the sum in Eq. (12) as leading approximation, as

$$V(\phi)\approx\frac{9}{(2\pi)^6R^4}\left[1-\cos\left(\frac{\phi}{f_{\rm eff}}\right)\right]. \tag{13}$$

1) If $1/R$ is $\mathcal{O}$(TeV), the theory could be relevant for the Higgs mechanism through the Hosotani mechanism [11, 12].

As can be seen in Table 1, this approximation is appropriate. Thus, the potential is identical to that of the natural inflation, given by Eq. (1), and the analytic calculations are straightforward, especially when ϕ is close to the top [14]. Here, we just give the results of the observable quantities: the power spectrum of the curvature perturbation, $\mathcal{P}_{\mathcal{R}}$, the corresponding spectral index, $n_{\mathcal{R}}$, the tensor-to-scalar ratio, r, and the non-linear parameter $f_{\rm NL}$ [15]. Under the slow-roll approximation, these parameters are given by

$$\mathcal{P}_{\mathcal{R}}^{1/2}=\sqrt{\frac{8V}{3\epsilon M_{\rm P}^4}}, \tag{14}$$

$$n_{\mathcal{R}}=1-6\epsilon+2\eta, \tag{15}$$

$$r=16\epsilon, \tag{16}$$

$$f_{\rm NL}=\frac{5}{6}(3\epsilon-\eta). \tag{17}$$

Here, ϵ and η are the usual slow-roll parameters defined by

$$\epsilon\equiv\frac{M_{\rm P}^2}{16\pi}\left(\frac{V'}{V}\right)^2, \tag{18}$$

$$\eta\equiv\frac{M_{\rm P}^2}{8\pi}\frac{V''}{V}, \tag{19}$$

where prime denotes a differential with respect to ϕ. it should be noted, that the running of $n_{\mathcal{R}}$, which in the slow-roll approximation can be written as

$$\frac{{\rm d}n_{\mathcal{R}}}{{\rm d}\log k}=-16\epsilon\eta+24\epsilon^2+2\xi^2, \tag{20}$$

where

$$\xi^2\equiv\frac{M_{\rm P}^4}{64\pi^2}\frac{V'V'''}{V^2} \tag{21}$$

is another slow-roll parameter, is second order in the slow-roll approximation and it is negligibly small compared with the other quantities[1)]; thus, the calculation is straightforward and not presented here. Also, the running of r [17],

$$\frac{{\rm d}\log r}{{\rm d}\log k}=2(2\epsilon-\eta), \tag{22}$$

is another first order quantity; thus, it is an observable. However, the same results can be obtained by combining Eqs. (15) and (16), which can serve as a consistency check. Writing Eqs. (14), (15), (16), and (17) in terms of $f_{\rm eff}$ and Rm such as in Eq. (12), we can obtain

$$\mathcal{P}_{\mathcal{R}}^{1/2}=\frac{8\sqrt{3}}{(2\pi)^{5/2}}\frac{f_{\rm eff}/M_{\rm P}}{(RM_{\rm P})^2}\left\{2-\frac{32\pi(f_{\rm eff}/M_{\rm P})^2}{16\pi(f_{\rm eff}/M_{\rm P})^2+1}\exp\left[\frac{-N}{8\pi(f_{\rm eff}/M_{\rm P})^2}\right]\right\}\left\{\frac{32\pi(f_{\rm eff}/M_{\rm P})^2}{16\pi(f_{\rm eff}/M_{\rm P})^2+1}\exp\left[\frac{-N}{8\pi(f_{\rm eff}/M_{\rm P})^2}\right]\right\}^{1/2}, \tag{23}$$

$$n_{\mathcal{R}}=1-\frac{1}{8\pi(f_{\rm eff}/M_{\rm P})^2}\left\{2+\frac{32\pi(f_{\rm eff}/M_{\rm P})^2}{16\pi(f_{\rm eff}/M_{\rm P})^2+1}\exp\left[\frac{-N}{8\pi(f_{\rm eff}/M_{\rm P})^2}\right]\right\}\left\{2-\frac{32\pi(f_{\rm eff}/M_{\rm P})^2}{16\pi(f_{\rm eff}/M_{\rm P})^2+1}\exp\left[\frac{-N}{8\pi(f_{\rm eff}/M_{\rm P})^2}\right]\right\}^{-1}, \tag{24}$$

$$r=\frac{1}{\pi(f_{\rm eff}/M_{\rm P})^2}\frac{32\pi(f_{\rm eff}/M_{\rm P})^2}{16\pi(f_{\rm eff}/M_{\rm P})^2+1}\exp\left[\frac{-N}{8\pi(f_{\rm eff}/M_{\rm P})^2}\right]\left\{2-\frac{32\pi(f_{\rm eff}/M_{\rm P})^2}{16\pi(f_{\rm eff}/M_{\rm P})^2+1}\exp\left[\frac{-N}{8\pi(f_{\rm eff}/M_{\rm P})^2}\right]\right\}^{-1}, \tag{25}$$

$$f_{\rm NL}=\frac{5}{48\pi(f_{\rm eff}/M_{\rm P})^2}\left\{1+\frac{16\pi(f_{\rm eff}/M_{\rm P})^2}{16\pi(f_{\rm eff}/M_{\rm P})^2+1}\exp\left[\frac{-N}{8\pi(f_{\rm eff}/M_{\rm P})^2}\right]\right\}\left\{2-\frac{32\pi(f_{\rm eff}/M_{\rm P})^2}{16\pi(f_{\rm eff}/M_{\rm P})^2+1}\exp\left[\frac{-N}{8\pi(f_{\rm eff}/M_{\rm P})^2}\right]\right\}^{-1}, \tag{26}$$

where

$$N\equiv\int H{\rm d}t \tag{27}$$

is the number of e-folds. Using Eq. (10), we can write these in terms of g_4 and R as

$$\mathcal{P}_{\mathcal{R}}^{1/2}=\frac{8\sqrt{3}}{(2\pi)^{7/2}g_4(RM_{\rm P})^3}\left\{2-\frac{8}{\pi(g_4RM_{\rm P})^2+4}\exp\left[-N\frac{\pi}{2}(g_4RM_{\rm P})^2\right]\right\}\left\{\frac{8}{\pi(g_4RM_{\rm P})^2+4}\exp\left[-N\frac{\pi}{2}(g_4RM_{\rm P})^2\right]\right\}^{1/2}, \tag{28}$$

$$n_{\mathcal{R}}=1-\frac{\pi}{2}(g_4RM_{\rm P})^2\left\{2+\frac{8}{\pi(g_4RM_{\rm P})^2+4}\exp\left[-N\frac{\pi}{2}(g_4RM_{\rm P})^2\right]\right\}\left\{2-\frac{8}{\pi(g_4RM_{\rm P})^2+4}\exp\left[-N\frac{\pi}{2}(g_4RM_{\rm P})^2\right]\right\}^{-1}, \tag{29}$$

1) In more general classes of inflation models [16], a sufficiently large ${\rm d}n_{\mathcal{R}}/{\rm d}\log k$ can be obtained.

$$r=\frac{32\pi(g_4RM_{\rm P})^2}{\pi(g_4RM_{\rm P})^2+4}\exp\left[-N\frac{\pi}{2}(g_4RM_{\rm P})^2\right]\left\{2-\frac{8}{\pi(g_4RM_{\rm P})^2+4}\exp\left[-N\frac{\pi}{2}(g_4RM_{\rm P})^2\right]\right\}^{-1}, \tag{30}$$

$$f_{\rm NL}=\frac{5}{12}\pi(g_4RM_{\rm P})^2\left\{1+\frac{4}{\pi(g_4RM_{\rm P})^2+4}\exp\left[-N\frac{\pi}{2}(g_4RM_{\rm P})^2\right]\right\}\left\{2-\frac{8}{\pi(g_4RM_{\rm P})^2+4}\exp\left[-N\frac{\pi}{2}(g_4RM_{\rm P})^2\right]\right\}^{-1}. \tag{31}$$

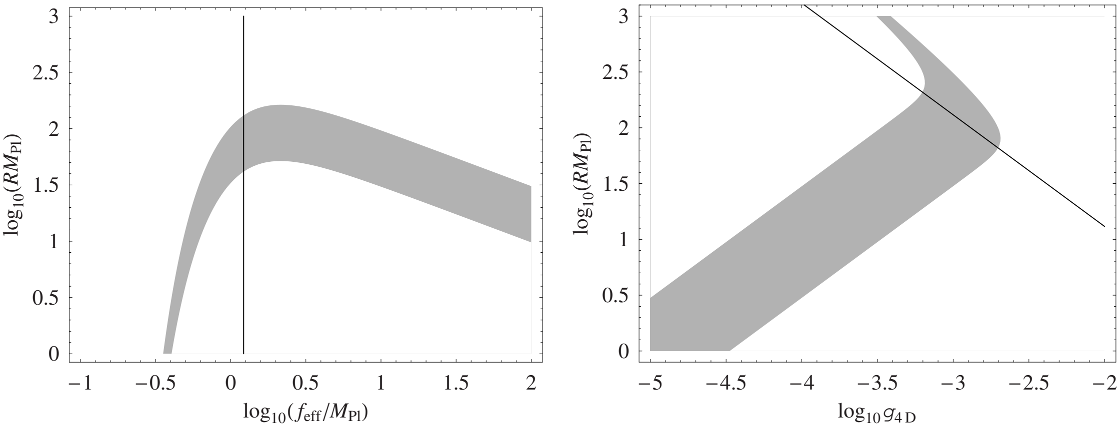

Fig. 1. Plots of $\mathcal{P}_{\mathcal{R}}^{1/2}$ on the $f_{\rm eff}/M_{\rm P}$-$RM_{\rm P}$ plane (left pane) and on the g_4-$RM_{\rm P}$ plane (right pane) determined at $N=60$. The shaded regions denote $10^{-5}\lesssim\mathcal{P}_{\mathcal{R}}^{1/2}\lesssim10^{-4}$, and the solid lines correspond to $n_{\mathcal{R}}=0.96$. It should be noted, that while for a large region $n_{\mathcal{R}}$ is saturated at $n_{\mathcal{R}}\approx0.967$ [see Eq. (32)], only a limited region is allowed for $\mathcal{P}_{\mathcal{R}}^{1/2}$.

Figure 1 shows $\mathcal{P}_{\mathcal{R}}^{1/2}$ and $n_{\mathcal{R}}$ determined at $N=60$ as functions of $f_{\rm eff}$, R, and g_4. We also compare the analytic estimations with numerical results in Table 1. As can be seen in Table 1, Eq. (13) is an appropriate approximation.

It can be seen from Eqs. (24), (25), and (26), that $n_{\mathcal{R}}$, r and $f_{\rm NL}$ depend only on the effective decay constant, $f_{\rm eff}$. This leads to the following simple expressions in the limit $f_{\rm eff}/M_{\rm P}\to\infty$, which is suitable for sufficiently long inflation[1)], as

$$n_{\mathcal{R}}\approx1-\frac{4}{1+2N}, \tag{32}$$

$$r\approx\frac{16}{1+2N}, \tag{33}$$

$$f_{\rm NL}\approx\frac{5}{3(1+2N)}. \tag{34}$$

Thus, it can be seen that in this limit, determined at certain e-folds before the end of inflation, these have definite values independent of $f_{\rm eff}$ or R. This is not unexpected, as the huge $f_{\rm eff}$ value indicates that the total number of e-folds we can obtain is very high, and the last 60 e-folds are only a final small fraction of the whole expansion. Therefore, in such a case, the physical properties become completely insensitive to the details of the model, as the inflationary dynamics is already following the late-time attractor. Thus, we obtain nearly identical values of $n_{\mathcal{R}}$, r, and $f_{\rm NL}$ in the limit $f_{\rm eff}/M_{\rm P}\to\infty$. This also indicates that the shapes of $\mathcal{P}_{\mathcal{R}}$ are identical, while only its overall amplitude depends on the inflationary energy scale[2)]

Figure 2 shows the r–$n_{\mathcal{R}}$ plot. It should be noted, that as shown in Eqs. (32) and (33), they are saturated as $f_{\rm eff}/M_{\rm P}\to\infty$, which corresponds to the upper right end of the curve where $n_{\mathcal{R}}\approx0.967$ and $r\approx0.132$. The shaded region shows the current observational 1σ bound $n_{\mathcal{R}}=0.960^{+0.014}_{-0.013}$ which is derived from the Planck data combined with the observations of type Ia supernovae (SN) and baryon acoustic oscillations (BAO) [5], and the points on the curve explicitly denote several constraints on $n_{\mathcal{R}}$: the central value $n_{\mathcal{R}}=0.960$ (circle) and the lower bound $n_{\mathcal{R}}^{\rm lower}=0.947$ (triangle). Our model is well below the upper bound $n_{\mathcal{R}}^{\rm upper}=0.974$ and there is no solution corresponding to this point (square). The corresponding values of r for $n_{\mathcal{R}}=0.960$ and $n_{\mathcal{R}}=0.947$ are 0.0528 and 0.0230, respectively. The current upper limit, $r<0.07$ (95% confidence level) [20], includes the wide predicted range of r of our model. For the observationally allowed range of $n_{\mathcal{R}}$, $0.01\lesssim r\lesssim0.07$, which is sufficiently large to be detected within a few years by the forthcoming cosmological experiments; therefore, it can serve as the

1) In the limit $f_{\rm eff}/M_{\rm P}\to\infty$, i.e., $g_4M_{\rm P}\ll1/(2\pi R)$, the gravitational force, which scales as $(m^2/M_*^2)/r^{2+n}$, where M_* is the cutoff mass scale in $4+n$ dimensions, becomes stronger than the gauge force between two Kaluza–Klein particles, g^2/r^{2+n}. In this parameter regime, the gravitational effects cannot be neglected and the effective potential can be modified: in this sense, the naive idea of extranatural inflation is as unnatural as that of natural inflation. See Ref. [18] for a more detailed discussion.

2) See, e.g., Fig. 3 in Ref. [19].

Table 1. Comparison of analytic estimations with numerical results. In the top row, R is chosen to transform the inflationary energy scale as $\Lambda = 10^{-3}M_{\rm P}$, $10^{-5/2}M_{\rm P}$, and $10^{-2}M_{\rm P}$. It should be noted that r is rather close to the observational sensitivity of the experiments planned in the near future. As can be seen, the leading approximation applying the $n=1$ part of Eq. (12) is reasonably good.

		$\mathcal{P}_{\mathcal{R}}^{1/2}$	$n_{\mathcal{R}}$	r
$\log_{10}(f_{\rm eff}/M_{\rm P})=0.00$	analytic	4.96×10^{-5}	0.952	0.032
$\log_{10}(RM_{\rm P})=2.04$	numerical	4.84×10^{-5}	0.955	0.033
$\log_{10}(f_{\rm eff}/M_{\rm P})=0.50$	analytic	1.25×10^{-5}	0.967	0.117
$\log_{10}(RM_{\rm P})=2.04$	numerical	1.33×10^{-5}	0.967	0.112
$\log_{10}(f_{\rm eff}/M_{\rm P})=1.00$	analytic	3.94×10^{-5}	0.967	0.131
$\log_{10}(RM_{\rm P})=1.54$	numerical	4.25×10^{-5}	0.967	0.130
$\log_{10}(f_{\rm eff}/M_{\rm P})=1.50$	analytic	1.25×10^{-5}	0.967	0.131
$\log_{10}(RM_{\rm P})=1.54$	numerical	1.33×10^{-5}	0.967	0.112
$\log_{10}(f_{\rm eff}/M_{\rm P})=2.00$	analytic	3.94×10^{-5}	0.967	0.132
$\log_{10}(RM_{\rm P})=1.04$	numerical	4.26×10^{-5}	0.967	0.134

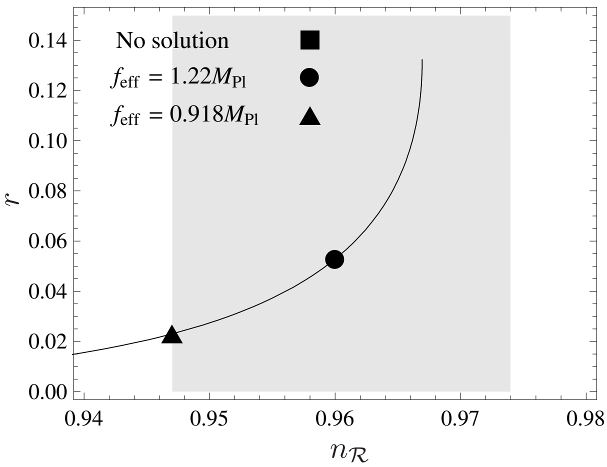

Fig. 2. Prediction of the model on the r–$n_{\mathcal{R}}$ plane. Both r and $n_{\mathcal{R}}$ are determined at 60 e-folds before the end of the inflation. The shaded region shows the current observational 1σ bound determined from the Planck, SN, and BAO data: $n_{\mathcal{R}} = 0.960^{+0.014}_{-0.013}$. Our model is well within the observational upper bound and has a wide parameter space satisfying the lower bound.

first observational test. It should also be noted that $f_{\rm NL}$ is always significantly smaller than 1; thus, the non-Gaussian signature is not observable at all.

After the end of the inflation, the inflaton starts oscillating at the global minimum. Although we assume no direct coupling between the hidden and the visible sectors, they can communicate gravitationally and the energy stored in the inflaton field can be converted to the light relativistic particles of the standard model to reheat the universe. The reheating temperature $T_{\rm RH}$ by the gravitational interaction can be determined in terms of the parameters of our model. With an interaction rate

$$\Gamma_{\rm grav} \sim \frac{m_\phi^3}{M_{\rm P}^2}, \tag{35}$$

using

$$m_\phi^2 \sim V'' \sim \frac{1}{f_{\rm eff}^2 R^4} = \frac{M_{\rm P}^2}{(f_{\rm eff}/M_{\rm P})^2 (RM_{\rm P})^4}, \tag{36}$$

we can rewrite Eq. (35) as

$$\Gamma_{\rm grav} \sim \frac{M_{\rm P}}{(f_{\rm eff}/M_{\rm P})^3 (RM_{\rm P})^6}. \tag{37}$$

Considering that inflation ends when $\dot{\phi}_{\rm end}^2 = V_{\rm end}$, we can find the Hubble parameter at the end of inflation, under the approximation in Eq. (13), as

$$\begin{aligned} H_{\rm end} &= \frac{3}{(2\pi)^{3/2}\pi}(RM_{\rm P})^{-1}[16\pi(f_{\rm eff}/M_{\rm P})^2+1]^{-1/2}R^{-1} \\ &\sim \mathcal{O}(0.1)\frac{R^{-1}}{(f_{\rm eff}/M_{\rm P})RM_{\rm P}}. \end{aligned} \tag{38}$$

Thus, for most of the parameter space $H_{\rm end} \gg \Gamma_{\rm grav}$ and the energy transfer occurs at a long time after the inflation. It can be seen, that the reheating temperature $T_{\rm RH}$ can be determined as [21]:

$$T_{\rm RH} \lesssim \mathcal{O}(0.1)\sqrt{\Gamma_{\rm grav} M_{\rm P}} \sim \mathcal{O}(0.1)\frac{M_{\rm P}}{(f_{\rm eff}/M_{\rm P})^{3/2}(RM_{\rm P})^3}. \tag{39}$$

As an example, if we apply $f_{\rm eff}/M_{\rm P}=1$ and $RM_{\rm P}=100$, the maximum reheating temperature is determined as $T_{\rm RH} \sim 10^{12-13}$GeV. Then, the Universe follows the well-known hot Big Bang evolution.

4 Conclusions

In this paper, we have presented a cosmological scenario of the hidden sector SU(2) gauge symmetry in a five-dimensional orbifold $M_4\times S^1/\mathbb{Z}_2$. The model is minimal in several aspects: a non-Abelian SU(2) gauge group of rank 1 is chosen and the minimal orbifold is considered as the extra dimension. Owing to the non-Abelian nature, the inflaton potential is fully radiatively generated, without introducing any ad hoc interactions or additional matter fields. Interestingly, the fully radiatively generated one-loop potential can support a sufficiently long period of slow-roll inflation within the theoretically desired parameter range where the theory is weakly coupled, i.e., $g_4 \ll 1$, during the inflationary epoch. The model predicts the observable cosmological quantities, which are in good agreement with the latest cosmological observations:

$$1.2\times10^{-5} \lesssim \mathcal{P}_{\mathcal{R}} \lesssim 4.9\times10^{-5}, \tag{40}$$

$$0.952 \lesssim n_{\mathcal{R}} \lesssim 0.966, \tag{41}$$

$$0.03 \lesssim r \lesssim 0.13. \tag{42}$$

The power spectrum of the curvature perturbation $\mathcal{P}_{\mathcal{R}}$ and the corresponding spectral index $n_{\mathcal{R}}$ are in good agreement with the current observations, while $f_{\rm NL}$ is always significantly smaller than 1 and no detectable non-Gaussianity is expected. Furthermore, the predicted tensor-to-scalar ratio r is quite close to the sensitivity of the cosmological experiments planned in near future, which can be the first test of our minimal cosmological model. Finally, the reheating temperature $T_{\rm RH}$ is determined to be sufficiently high to successfully follow the standard hot Big Bang evolution.

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