Inflation from Broken Scale Invariance

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Inflation from Broken Scale Invariance

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We construct a model of inflation based on a low-energy effective theory of spontaneously broken global scale invariance. This provides a shift symmetry that protects the inflaton potential from quantum corrections. Since the underlying scale invariance is non-compact, arbitrarily large inflaton field displacements are readily allowed in the low-energy effective theory. A weak breaking of scale invariance by almost marginal operators provides a non-trivial inflaton minimum, which sets and stabilizes the final low-energy value of the Planck scale. The underlying scale invariance ensures that the slow-roll approximation remains valid over large inflaton displacements, and yields a scale invariant spectrum of perturbations as required by the CMB observations.

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Inflation is the leading contender for the explanation of why the Universe is so big, old, and smooth [1–3]. It also predicts the initial spectrum of almost scale invariant density fluctuations [4]. These inflationary fluctuations excellently fit the cosmic microwave background (CMB) measurements by WMAP and Planck. Very recently the BICEP2 experiment claimed an observation of CMB polarization [5], which fit the spectrum of primordial gravity waves [6] that can also be created during inflation. The BICEP2 results, if due to primordial gravity waves, point towards large field models of inflation, to explain the claimed large tensor-to-scalar ratio. Such models involve large field changes $\Delta\varphi > M_{Pl}$ during inflation, and need a very flat and small potential in Planck units [3]. They are difficult to realize because at large field values the quantum corrections can be large. However, setups using a pseudo-Goldstone boson of some weakly broken symmetry as the inflaton [7, 8], have an approximate shift symmetry which protects the potential from large corrections [9–11]. The inflaton's shift symmetry is a 'phase rotation', and the inflaton is necessarily a pseudo-scalar (essentially a type of axion). Here we argue that an alternative is to use a scalar Goldstone boson for a non-compact, spontaneously broken global scale symmetry, the dilaton, as the inflaton. This automatically accommodates large field variations since the symmetry and the vacuum manifold are non-compact. Scale invariance forbids a direct Einstein-Hilbert term in the action, so the leading operator controlling graviton dynamics is a dilaton-graviton coupling $\Phi^2 R$. The Planck scale arises from the dilaton VEV $\langle\Phi\rangle \sim M_{Pl}$. A fully scale invariant theory allows only a quartic dilaton self coupling, without a non-trivial minimum, protected from loop corrections by an effective shift symmetry which arises from the underlying scale symmetry. An inclusion of small explicit breaking terms yields a non-trivial dilaton VEV at large but finite values $O(M_{Pl})$ with a very flat potential. All corrections to the inflaton potential will be suppressed by the small parameters characterizing the sizes of the explicit breaking terms.

Our main assumption is that the low-energy effective Lagrangian is approximately scale invariant. Global scale transformations are given by $x^\mu \to \bar{x}^\mu = e^{-\lambda} x^\mu$, or equivalently $g_{\mu\nu} \to e^{-2\lambda} g_{\mu\nu}$. These have the effect $R \to e^{2\lambda} R$ on the scalar curvature, while generic operators transform as $\mathcal{O} \to e^{\lambda\Delta}\mathcal{O}$, where Δ is the scaling dimension of $\mathcal{O}$. The spontaneous breaking of scale invariance is parameterized by the dilaton field Φ, which is the Goldstone boson for broken scale invariance, and which is the inflaton in our setup. Once the dilaton is stabilized by the small explicit breaking terms, its VEV will give rise to the effective Planck scale. We will assume that initially the dilaton is displaced far from its minimum, and that its rolling to its minimum drives inflation.

The general scale invariant Lagrangian that we will be considering is given by

$$\begin{aligned} \mathcal{L} \;=\; & \sqrt{-g}\left[\tilde{\xi}\Phi^2 R - \frac{1}{2}(\nabla\Phi)^2 - V(\Phi)\right] + \Delta\mathcal{L}(g_{\mu\nu}, \Phi) \\ + \; & \mathcal{L}_M(g_{\mu\nu}, \Phi, \Psi) \; . \end{aligned} \tag{1}$$

where R is the Ricci scalar and $\tilde{\xi}$ is a dimensionless parameter. Note, that scale invariance forbids the presence of the usual Einstein-Hilbert term. The potential $V(\Phi)$ will be specified below, but exact scale invariance would require $V(\Phi) = \alpha^2\Phi^4$, with a constant α. Scale invariance forbids large corrections to the dilaton potential, hence eliminating the η-problem. This remains valid even after including the loop corrections from the interactions with other fields, as long as these fields do not violate scale invariance explicitly. This will be the case if the masses of the fields interacting with the dilaton originate from the dilaton VEV itself, in which case the resulting corrections will just renormalize the coefficient of the Φ^4 coupling. In order to recover Einstein gravity, the potential must give rise to a non-vanishing VEV for Φ, $\langle\Phi\rangle^2 = M_{Pl}^2/2\tilde{\xi}$. This requires the presence of small explicit breaking terms, whose corrections to the dilaton potential will nevertheless be suppressed by the small parameter characterizing the magnitude of the explicit breaking. This follows since the theory - including the regulator - has a manifest (non-linearly realized) shift symmetry, which arises from scale invariance after field redefinitions. This also guarantees that all the perturbative graviton loop corrections are completely under control, much like in the case of axion monodromy [11] (see also [12]). $\Delta\mathcal{L}(g_{\mu\nu},\Phi)$ contains operators with extra derivatives and inverse powers of Φ, for example the Weyl term involving R^2 would be in this part of the Lagrangian. $\mathcal{L}_M(g_{\mu\nu},\Phi,\Psi)$ contains any other dynamics involving fields collectively denoted by Ψ (such as the Standard Model (SM) fields), which may or may not be coupled to Φ (but they certainly couple to the metric in order to preserve Lorentz invariance). We will discuss the role of these two terms later.

In order to understand the inflationary dynamics of this system, it is convenient to perform a Weyl transformation of the metric and go to the Einstein frame $g_{\mu\nu} \to \Omega^2 g_{\mu\nu}$, where $\Omega = \Omega(x)$ satisfies $\Omega^2\tilde{\xi}\Phi^2 = M_P^2/2$. The rescaled Lagrangian is given by

$$\mathcal{L} = \sqrt{-g}\left[\frac{M_{Pl}^2}{2}R - \frac{1}{2}(\nabla\varphi)^2 - V(\varphi)\right] + \Delta\mathcal{L}\big(\Omega^2(\varphi)g_{\mu\nu},\Phi(\varphi)\big) + \mathcal{L}_M\left(\Omega^2(\varphi)g_{\mu\nu},\Phi(\varphi),\Psi\right) \quad (2)$$

where $V(\varphi) = \frac{M_{Pl}^4}{4\tilde{\xi}^2}\frac{V(\Phi(\varphi))}{\Phi^4(\varphi)}$. The relation between the original dilaton and the Einstein frame inflaton φ is given by (with boundary condition $\Phi(\varphi=0) = \langle\Phi\rangle$)

$$\Phi(\varphi) = \langle\Phi\rangle \exp\left(\frac{\sqrt{\xi}\varphi}{M_{Pl}}\right) \ , \quad \frac{1}{\xi} = \frac{1}{2\tilde{\xi}} + 6 \ . \quad (3)$$

In this frame the original scale invariance of the theory will manifest itself in a shift symmetry for the inflaton

$$\varphi \to \bar{\varphi} = \varphi + \frac{M_{Pl}}{\sqrt{\xi}}\lambda \ . \quad (4)$$

Thus Eq. (2) can be thought of as the non-linearly realized Lagrangian for the spontaneously broken non-compact group of scale transformations, where the above shift symmetry is the remnant of the original scale invariance. The Einstein-Hilbert term is shift symmetric, since it does not contain φ. The kinetic term for the scalar is shift symmetric because it contains only derivatives. The scalar potential term $V(\varphi)$ becomes a constant (if we started out with a quartic Φ^4 in the Jordan frame, as required in the absence of explicit breaking terms). The terms in $\Delta\mathcal{L}$ already contain derivatives of φ only, and thus will be obviously shift invariant. The only non-trivial terms are those that involve matter fields coupled to φ in $\mathcal{L}_M$: here explicit powers of $e^{\sqrt{\xi}\varphi/M_{Pl}}$ will appear from the Weyl transformation of the metric, seemingly giving rise to non-derivative interactions. The important point is that such factors will also be present in the kinetic terms of the matter fields: once the matter fields are suitably redefined in order to canonically normalize their kinetic terms, the inflaton will again appear only derivatively coupled, obeying the shift symmetry. Hence all the terms in Eq. (2) which were originally exactly scale invariant remain invariant under the shift symmetry.

Notice also that, given $\varphi = (M_{Pl}/\sqrt{\xi})\log(\Phi/\langle\Phi\rangle)$, if the dilaton field starts out at small values $\Phi_0 \sim 0$ far from the minimum of the potential and moves out to $\langle\Phi\rangle \sim M_{Pl}$, the field space range for φ can be larger than M_{Pl} without ever leaving the regime of validity of the effective theory. For example assuming $\Phi_0 \sim 10^{-15}\langle\Phi\rangle \sim$ TeV, we find $|\Delta\varphi| \sim 15M_{Pl}$, a seemingly super-Planckian field excursion in the Einstein frame.

The scale invariant $\alpha^2\Phi^4$ dilaton potential yields a completely flat constant potential independent of φ in the Einstein frame. This is again a consequence of the shift symmetry Eq. (4). However for a completely flat potential the VEV $\langle\Phi\rangle$ (and the Planck scale) remain undetermined. One needs to systematically incorporate small explicit breaking terms into the Lagrangian which can fix the dilaton VEV at large values. Such explicit breaking terms could possibly originate from the interactions with additional matter contained in $\mathcal{L}_M$, in particular they could potentially be due to interactions with the SM fields. As long as the explicit breaking induced by these terms is weak, the shift symmetry Eq. (4) will remain approximately valid, and will continue to protect the low energy theory Eq. (2) from large corrections. We now consider several simple but well-motivated forms of potentials that systematically incorporate small explicit breakings of scale invariance, with a vanishing cosmological constant at the minimum (notice that scale invariance does not in itself say anything about the cosmological constant [13]). More examples which yield arbitrary power-law inflaton potentials protected by approximate shift symmetry are discussed in [14].

The first example takes the effect of a single marginally relevant operator with dimension $4-\epsilon$ into account. This type of potential [15, 16] naturally shows up in warped extra dimensions [17] after modulus stabilization via the Goldberger-Wise mechanism [18, 19] (which indeed corresponds to turning on a marginally relevant operator in the dual conformal field theory language). The resulting approximately scale invariant potential is

$$V(\Phi) = \Phi^4\left(\alpha + \beta\Phi^{-\epsilon}\right)^2 \ , \quad (5)$$

where ϵ corresponds to the anomalous dimension of the operator breaking scale invariance, $\epsilon \ll 1$. This potential is minimized at $\langle\Phi\rangle = (-\alpha/\beta)^{1/\epsilon}$, where it vanishes to reproduce an (approximately) zero vacuum energy density at the end of inflation. The inflaton potential in the Einstein frame reads

$$V(\varphi) = \frac{M_{Pl}^4}{4}\frac{\alpha^2}{\tilde{\xi}^2}\left(1 - e^{-\epsilon\sqrt{\xi}\varphi/M_{Pl}}\right)^2 \ . \quad (6)$$

This is a very flat potential, as long as $\epsilon \ll 1$: a result of the small explicit breaking of scale invariance. Note that the form of the potential Eq. (6) is the same as that of the Starobinsky model [20], with the important difference that the exponent here is controlled by the amount of explicit breaking in the field theory. In contrast, in the original Starobinsky model the exponent is fixed by $4D$ general covariance. To understand why the Starobinsky potential is a special case of Eq. (6), however, all one needs is scaling symmetry. The starting action of [20] can be thought of as a special case of scale invariant theory where the breaking of scale invariance is induced purely gravitationally, by an explicit $M_{Pl}^2 R$ term. This immediately explains the necessity that in Starobinsky inflation, the R^2 term *must* dominate over $M_{Pl}^2 R$ to yield inflation: the scale symmetry breaking term *must* be subleading in the UV for the protection mechanism to be operational. This is also the reason behind the emergence of the same type of potentials in the context of induced gravity, as explained in [21].

The slow-roll parameters and the number of e-folds of inflation are given in this model by

$$\epsilon_V = \frac{2\epsilon^2\xi}{(1-e^{\epsilon\sqrt{\xi}\varphi/M_{Pl}})^2}\ , \quad \eta_V = \epsilon_V\left(2 - e^{\epsilon\sqrt{\xi}\varphi/M_{Pl}}\right)\ ,$$
$$N \simeq \frac{1}{2\epsilon^2\xi}\left[\left(e^{\epsilon\sqrt{\xi}\varphi_0/M_{Pl}} - 1\right) - \frac{\varphi_0}{\sqrt{2}M_{Pl}}\right]\ . \tag{7}$$

The above expressions depend only on the combination $\epsilon\sqrt{\xi}$, which is the single parameter needed to characterize this model. Now we can compute the scalar power spectrum $\mathcal{P}_s$, the tensor-to-scalar ratio $r \simeq 16\epsilon_V$, and the tilt of the primordial scalar perturbations $n_s \simeq 1 + 2\eta_V - 6\epsilon_V$, at CMB horizon exit with $N_{cmb} \simeq 60$. We show in Fig. 1 the values of n_s and r while varying $\epsilon\sqrt{\xi} \in [-0.5, 0.5]$, for $\varphi_0 < \langle\varphi\rangle = 0$, corresponding to almost marginal perturbations. The same results are obtained for $\varphi_0 > \langle\varphi\rangle = 0$, but with opposite signs for ϵ. The points shown correspond to $\epsilon\sqrt{\xi} = -0.001, -0.01, -0.05, 0.001, 0.01, 0.1, 0, 5$, and $\sqrt{2/3}$ corresponding to the Starobinsky model. If we insist on solutions with $\varphi_0 < \langle\varphi\rangle$, we can see that this model can accommodate both very small values of r (for relatively large anomalous dimensions $\epsilon\sqrt{\xi} \sim O(0.1)$, that is marginally relevant perturbations), while r can be pushed into the region favored by BICEP2 for $\epsilon\sqrt{\xi} < 0$, corresponding to marginally irrelevant perturbations. Similar observations were noted in specific constructions in [22]. The COBE normalization $(\mathcal{P}_s)_{exp} \sim 10^{-9}$ enforces a constraint on the parameter α in the potential, for fixed ϵ and $\tilde{\xi}$. Explicitly one obtains

$$\mathcal{P}_s = \frac{\alpha^2}{24\pi^2\tilde{\xi}^2}\frac{\sinh^4(\epsilon\sqrt{\xi}\varphi_{cmb}/2M_{Pl})}{\epsilon^2\xi}\ , \tag{8}$$

where φ_{cmb} is a function of $\epsilon\sqrt{\xi}$. Since $\mathcal{P}_s$ increases with $\epsilon\sqrt{\xi}$, smaller values of the explicit breaking parameter ϵ — and therefore better slow-roll approximation — accommodate the observed power spectrum more easily, as

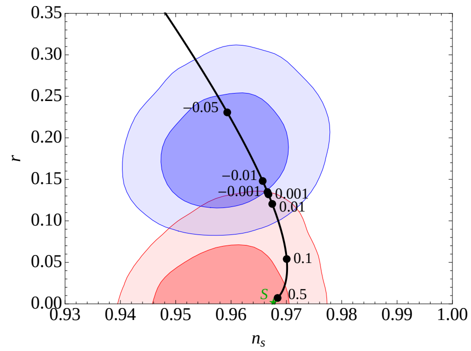

FIG. 1: Values of n_s and r for $\epsilon\sqrt{\xi} \in [-0.5, 0.5]$, for $\varphi_0 < \langle\varphi\rangle = 0$. The same results are obtained for $\varphi_0 > \langle\varphi\rangle = 0$, but with opposite signs for ϵ. The points shown correspond to $\epsilon\sqrt{\xi} = -0.001, -0.01, -0.05, 0.001, 0.01, 0.1, 0, 5$, and $\sqrt{2/3}$ in green for the Starobinsky model. The red and blue contours show the 68% and 95% confidence regions by Planck and BICEP2 respectively.

expected from general inflationary phenomenology. From the minimum of Eq. (5) one naturally expects that ϵ is of the order of $1/\ln(M_{Pl}/\Lambda_\epsilon)$, where Λ_ϵ parametrizes the onset of scaling symmetry breaking. For instance $\Lambda_\epsilon \sim 10^{\pm 3}M_{Pl}$ yields $\epsilon \sim 0.1$, while $\Lambda_\epsilon \sim 10^{\pm 17}M_{Pl}$ gives $\epsilon \sim 0.01$. Using $\epsilon\sqrt{\xi} = \pm 0.01$ and for the most favorable case of $\tilde{\xi} \simeq 16\pi^2$ (notice that $\mathcal{P}_s$ decreases with increasing $\tilde{\xi}$), the scalar power spectrum is

$$\mathcal{P}_s \simeq \left(\frac{\alpha}{0.1}\right)^2 \times 10^{-9}\ , \tag{9}$$

which requires a perturbative value of α compared to its NDA estimate $\alpha \sim 4\pi$.

Another simple potential could arise in the presence of a marginally relevant and a marginally irrelevant perturbation. For simplicity we take their dimensions to be $4 \pm \epsilon$, though they could be independent. So,

$$V(\Phi) = -\alpha^2\Phi^4 + \beta^2\Phi^{4-\epsilon} + \gamma^2\Phi^{4+\epsilon}\ , \tag{10}$$

while in the Einstein frame

$$V(\varphi) = \frac{M_{Pl}^4}{4}\frac{\alpha^2}{\tilde{\xi}^2}\left(\cosh(\epsilon\sqrt{\xi}\varphi/M_{Pl}) - 1\right)\ . \tag{11}$$

This potential is clearly the non-compact analogue of the generic axion-type potentials for the case of a broken compact symmetry. Note, that the analogue of the axion decay constant appearing here is effectively given by $M_{Pl}/\epsilon\sqrt{\xi}$, which can be $\gg M_{Pl}$ for small ϵ. However, obtaining a 'large decay constant' and allowing for an even larger range of variation of φ is straightforward here. The cosmological parameters are

$$\epsilon_V = \frac{1}{2}\epsilon^2\xi\coth^2(\epsilon\sqrt{\xi}\varphi/2M_{Pl})\ , \quad \eta_V = \frac{\epsilon_V}{\cosh(\epsilon\sqrt{\xi}\varphi/M_{Pl})}\ ,$$
$$N \simeq \frac{2}{\epsilon^2\xi}\log\left[\cosh(\epsilon\sqrt{\xi}\varphi/2M_{Pl})\right]\ , \tag{12}$$

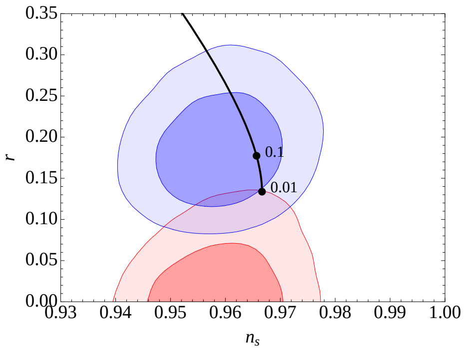

FIG. 2: Line of values of n_s and r for $\epsilon\sqrt{\xi} \in (0, 0.5]$, with points at $\epsilon\sqrt{\xi} = 0.1, 0.01$, for either sign of φ_0. The same results are obtained for negative ϵ. The red and blue contours show the 68% and 95% confidence regions by Planck and BICEP2 respectively.

which again only depend on the combination $\epsilon\sqrt{\xi}$. In Fig. 2 we show the line of values of n_s and r for $\epsilon\sqrt{\xi} \in (0, 0.5]$, with points at $\epsilon\sqrt{\xi} = 0.1, 0.01$, for either sign of φ_0. The same results are obtained for negative ϵ. Small values of $|\epsilon|$ yield approximately the same result as for $\epsilon = 0.01$ (which is also very similar to the result at small ϵ for the previous potential, see Fig. 1). Thus this particular model predicts a relatively large tensor-to-scalar ratio $r \gtrsim 0.1$. This is not surprising since the potential is an extrapolation of the quadratic potential, which generically yields larger r [3, 11]. The normalization of the scalar power spectrum is again approximately given by Eq. (9) for the same choice of parameters $\epsilon, \xi, \tilde{\xi}$.

Understanding the regime of validity of our effective field theory is straightforward in the Einstein frame where the inflaton 'decay constant', associated to the spontaneous breaking of scale invariance, is $f = M_{Pl}/\sqrt{\xi}$. The cutoff is at or below $\Lambda_{UV} = 4\pi M_{Pl}/\sqrt{\xi}$. We can explicitly check this by studying the operators at higher order in derivatives encoded in $\Delta\mathcal{L}$ in Eq. (1), and identifying the effective cutoff scale that suppresses them. One such term is R^2, which in the Einstein frame gives rise to

$$\frac{1}{g_R^2}R^2 \to \frac{1}{g_R^2}\left[R + 6\left(\frac{\sqrt{\xi}}{M_{Pl}}\nabla^2\varphi - \frac{\xi}{M_{Pl}^2}(\nabla\varphi)^2\right)\right]^2 .$$

Each of the terms on the r.h.s. indicates that the cutoff lies at, or somewhat below, Λ_{UV}. For instance, the R^2 term can be regarded as arising from integrating out a scalar of mass $M_R^2 \simeq g_R^2 M_{Pl}^2$, which for the NDA estimate $g_R \sim 4\pi$, sets the cutoff at $\Lambda_{UV} \approx M_R \sim 4\pi M_{Pl}$. Similarly, the other two terms set the cutoff at $\Lambda_{UV} \approx (g_R^2/\xi)M_{Pl} \sim 4\pi M_{Pl}$. Notice however that by taking small values of $\tilde{\xi}$ in Eq. (3), for which $\xi \simeq 2\tilde{\xi}$, this latter cutoff can be raised above the naive expectation, contrary to the R^2 case. The same behavior as for R^2 is found for the $R_{\mu\nu}^2/\tilde{g}_R^2$ operator. In this case it corresponds to a spin-2 ghost field with mass $\tilde{M}_R^2 \simeq \tilde{g}_R^2 M_{Pl}^2$. As long as $\tilde{g}_R^2$ is sufficiently large, the cutoff is above M_{Pl}.

We stress again that since the inflaton is derivatively coupled – it appears through its derivatives $\nabla\varphi$, in any of the operators in $\Delta\mathcal{L}$ – the field excursion of the inflaton beyond Λ_{UV} is not a problem given that the inflaton potential is almost flat. Large φ values could be problematic in non-derivative terms, associated with the explicit breaking of the shift symmetry. However – as long as the breaking of the scaling/shift symmetry is weak – they are small and under control, via ϵ-suppression. Even if the actual explicit breaking of scaling symmetry is below M_{Pl} but is weak, the low energy theory remains extremely well protected by the approximate shift symmetry, essentially staying valid all the way up to the scale of quantum gravity, because the scaling symmetry breaking sector is very efficiently sequestered from the low energy inflaton.

Finally we turn to the dynamics of the matter fields, which is clearly dependent on the UV completion. Assuming that at very high energies the SM fields are the proper degrees of freedom, the couplings in the matter Lagrangian $\mathcal{L}_M$ are classically marginal, with the exception of the Higgs mass. Thus at tree-level the SM Lagrangian is scale invariant, while the Higgs mass parameter constitutes a small explicit breaking of $O(m_H^2/M_{Pl}^2)$. At loop-level the SM couplings run, but the β-functions at high energies are perturbatively small, of $O(1/16\pi^2)$. The exact form of the couplings between the dilaton/inflaton and the SM matter fields depend on the details of the embedding of the SM fields into the scale invariant UV theory. To obtain their couplings, one can dress the dimensionful parameters with the appropriate powers of $\Phi/\langle\Phi\rangle = e^{\varphi/f}$, with $f = M_{Pl}/\sqrt{\xi}$. This leads to a decay rate to W, Z and h bosons of

$$\Gamma_\varphi \simeq \frac{4\xi}{32\pi}\frac{m_\varphi^3}{M_{Pl}^2} \simeq 0.5 \text{ GeV} \left(\frac{\xi}{1/12}\right)\left(\frac{m_\varphi}{10^{13}\text{ GeV}}\right)^3 , \tag{13}$$

where the mass of the inflaton, in the simplest example of Eq. (6) is given by

$$m_\varphi = M_{Pl}\frac{\alpha\epsilon\sqrt{\xi}}{\tilde{\xi}} \simeq 10^{13}\left(\frac{\alpha}{0.1}\right)\left(\frac{\epsilon\sqrt{\xi}}{0.01}\right)\left(\frac{16\pi^2}{\tilde{\xi}}\right) \text{ GeV} . \tag{14}$$

The reheat temperature is generically dominated by $\varphi \to WW, ZZ, hh$ decays, and is given by

$$T_{RH} \sim g_*^{-1/4}(\Gamma M_{Pl})^{1/2} \sim 3 \times 10^8 \text{ GeV} , \tag{15}$$

for $g_* \sim O(100)$ and for the parameters chosen above. We can see that this temperature is high enough to accommodate baryogenesis, but sufficiently low to avoid restoration of high scale symmetries (like GUT) and prevent any regeneration of undesired topological defects. So, in closing we note that our construction represents a fully viable model of large field inflation.

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