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Loop inflection-point inflation

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ABSTRACT

A novel inflection-point inflation model is analysed. The model considers a massless scalar field, whose self-coupling's running is stabilised by a non-renormalisable operator. The running is controlled by a fermion loop. We find that successful inflation is possible for a natural value of the Yukawa coupling $y \simeq 4 \times 10^{-4}$. The necessary fine-tuning is only $\sim 10^{-6}$, which improves on the typical tuning of inflection-point inflation models, such as MSSM inflation. The model predicts a spectral index within the 1-σ bound of the latest CMB observations, with a very small negative running, and negligible tensors ($r \sim 10^{-(9-10)}$). These results are largely independent of the order of the stabilising non-renormalisable operator.

1. Introduction

Cosmic inflation is an organic part of concordance cosmology. With a single stroke inflation addresses the fine-tuning problems of the hot big bang; namely the horizon and flatness problems and also produces the primordial curvature perturbation, which seeds structure formation and is in excellent agreement with CMB observations [1]. According to the inflationary paradigm, the Universe undergoes inflation when dominated by the potential density of a scalar field, called the inflaton. However, the identity of the inflaton is as yet unknown.

The latest CMB observations suggest that the scalar potential of the inflaton features an inflationary plateau (e.g. see Ref [2]). Numerous mechanisms have been put forward to generate such a plateau, involving exotic constructions in the context of elaborate, beyond-the-standard-model theories, such as superstrings. One such example is inflection-point inflation, where the inflationary plateau is due to the interplay of opposing contributions in the scalar potential, which (almost) cancel each other out generating a step on the otherwise steep potential wall. The original model was called A-term inflation, because it employed the A-term of a supersymmetric theory [3,4], or MSSM inflation, because it considered a flat direction in MSSM [5] as the inflaton. However, other models of inflection-point inflation have also been constructed [6,7]. Most of these also consider an elaborate setup in the context of supersymmetry, string theory or other extensions of the Standard Model.

However, an advantage of the idea of inflation is that it does not have to rely on exotic physics, in contrast to alternatives like the ekpyrotic scenario [8] or string gas cosmology [9]. Indeed, inflation may be realised simply within field theory in curved spacetime. It is also possible to achieve inflection-point inflation in this way. In this paper we explore such a possibility, where we exploit the loop corrections to the inflaton potential to generate the step-like plateau. This is similar to the works in Ref. [7]. However, in Ref. [7] the authors consider a rather complicated running of the inflaton self-coupling, where many particles are contributing to it. We consider a simpler setup.

In previous works care was taken so that loop corrections do not spoil the stability of the potential [10]. In contrast, here we consider a model in which the Coleman–Weinberg potential is unstable. Stability is recovered by introducing a Planck-suppressed effective operator.

We use natural units where $c = \hbar = 1$ and $8\pi G = m_P^{-2}$, with $m_P = 2.43 \times 10^{18}$ GeV being the reduced Planck mass.

2. Coleman–Weinberg Potential

The general expression for the 1-loop potential is given by the Coleman–Weinberg (CW) result [11]

$$V_{\rm eff} = V + \sum_{i=1}^{n} \frac{g_i M_i^4(\phi)}{64\pi^2} \ln\left(\frac{M_i^2(\phi)}{\mu^2}\right), \tag{1}$$

where V is the tree-level potential, μ is the renormalisation scale and M_i and g_i are, respectively, the field dependent tree level mass and the number of intrinsic degrees of freedom of the particle-i coupled with ϕ. We assume a quartic tree-level potential for the

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inflaton field

$$V = \lambda\,\phi^4\,, \tag{2}$$

and that the dominant contribution in Eq. (1) is given by the Yukawa coupling y between ϕ and a Weyl fermion.[1] Therefore we can approximate Eq. (1) with

$$V_{\rm eff}(\phi) = \left[\lambda - \beta\ln\left(\frac{y^2\phi^2}{\mu^2}\right)\right]\phi^4\,, \tag{3}$$

where we used Eq. (2) and $\beta = y^4/32\pi^2$. We can improve the potential by inserting the running expression for λ. Since we assumed that the Yukawa coupling y is the dominant contribution, a good approximation[2] for the RGE solution of λ is

$$\lambda(\mu) = \lambda(M) - 2\beta\log\left(\frac{\mu}{M}\right), \tag{4}$$

where M is the scale at which we impose the boundary condition on the running of λ. Since we are interested in studying a configuration in which the CW potential is unstable, it is natural to pick[3] $\lambda(M) = 0$. Using this and inserting Eq. (4) into Eq. (3) we get

$$V_{\rm eff}(\phi) = -\beta\ln\left(\frac{y^2\phi^2}{M^2}\right)\phi^4\,. \tag{5}$$

3. Inflation model with inflection point

The potential in Eq. (5) is not stable because it is unbounded from below. We assume that stability is ensured by the intervention of a non-renormalisable Planck-suppressed effective operator. Therefore let us consider the following inflaton potential

$$V = -\beta\ln\left(\frac{y^2\phi^2}{M^2}\right)\phi^4 + \lambda_n\frac{\phi^{2n+4}}{m_P^{2n}}, \tag{6}$$

where the first term is the 1-loop effective potential obtained in Eq. (5) and the second term is an effective non-renormalisable operator, with $\lambda_n \ll 1$ and $n \geq 1$. We consider only the dominant non-renormalisable term, of order n.

For the moment we choose $n = 1$ but later on we consider higher values of n. For simplicity, we study the model where

$$\frac{y^2}{M^2} = \frac{1}{m_P^2}. \tag{7}$$

If $y < 1$ (required for pertubativity), it is possible to realise such a condition with sub-Planckian M.

A priori, M and y can take whatever possible value. However it is possible to reduce the parameters space, identifying a preferred region which is essentially described by Eq. (7). For example, assuming that our inflaton is not the Higgs boson of the SM, it is reasonable to expect new physics to happen around the scale of grand unification (GUT-scale). Therefore it is reasonable to consider $M \sim 10^{15-16}$ GeV. In addition to that, the Yukawa coupling, y, generating the loop correction must be small enough to preserve perturbativity, but on the other side, also big enough to give rise to relevant corrections. Therefore a reasonable range for y is[4] around $10^{-(2-3)}$. Combining the two expected regions for M and y, we get that y/M is around $1/m_P$, therefore for the first analysis, in which we present a new idea for inflection point models, it is enough to study the model implementing Eq. (7). We will consider a broader range of M and y values in a future article.

Noting that the slow-roll formalism is independent of the potential normalisation, we reparameterise the potential as

$$V = \beta\left[-\ln\left(\frac{\phi^2}{m_P^2}\right)\phi^4 + \alpha\frac{\phi^6}{m_P^2}\right], \tag{8}$$

where $\alpha = \lambda_1/\beta$. Such a potential has a flat inflection point at

$$\phi_f = e^{1/4}m_P \quad\text{and}\quad \alpha_f \equiv \frac{2}{3\sqrt{e}}\,. \tag{9}$$

To study the inflationary predictions for values of α around α_f, we parameterise:

$$\alpha = (1+\delta)\alpha_f \tag{10}$$

and use δ as a free parameter. Varying δ allows us to find the range of allowed slopes of the plateau around the flat inflection point. Increasing δ increases the slope of the plateau. Decreasing δ to negative values introduces a local maximum.

There are two aspects to consider when constraining δ. First, by contrasting the computed inflationary observables with the observations. Second, by ensuring that the necessary remaining e-folds of inflation since the cosmological scales exited the horizon, N_*, is not greater than the total e-folds of inflation, $N_{\rm tot}$. When the parameter space for δ is established we calculate predictions for the inflationary observables, namely the spectral index of the scalar curvature perturbations, n_s, its running, $n_s' \equiv \frac{dn_s}{d\ln k}$ and the tensor-to-scalar ratio, r.

3.1. Computing N_*

First we must make clear the distinction between $N_{\rm tot}$ and N_*. $N_{\rm tot}$ depends mainly on the initial conditions of the inflaton. We set the beginning of inflation to be determined by $\epsilon = 1$, where $\epsilon = -\dot{H}/H^2$ is the usual slow-roll parameter. For the e-folds of *observable* inflation N_*, typically the reheating temperature has a large impact. However, our model does not need an in-depth investigation into reheating since in this model, after inflation, the field oscillates in a quartic minimum because of Eq. (2) and also

$$\lim_{\phi\to 0}\left[-\beta\ln\left(\frac{\phi^2}{m_P^2}\right)\phi^4\right] = \frac{1}{2}\beta\phi^4\,. \tag{11}$$

The average density of a scalar field coherently oscillating in a quartic potential scales as $\rho_\phi \propto a^{-4}$ [12], just as the density of a radiation dominated Universe. Hence, there is little distinction in the expansion between inflaton oscillations and radiation domination after reheating, which means that N_* is independent of the inflaton decay rate.

In this case we have

$$N_* = 62.8 - \ln\left(\frac{k}{a_0H_0}\right) + \frac{1}{3}\ln\left(\frac{g_*}{106.75}\right) + \frac{1}{3}\ln\left(\frac{V_{\rm end}^{1/4}}{10^{16}\rm GeV}\right) \tag{12}$$

where $k = 0.05\mathrm{Mpc}^{-1}$ is the pivot scale, $(a_0H_0)^{-1}$ is the comoving Hubble radius today, g_* is the effective number of relativistic degrees of freedom and $V_{\rm end} \equiv V(\phi_{\rm end})$, with 'end' denoting the end of inflation. This simplifies when we take $g_* = 106.75$, corresponding to the standard model at high energies. Inputting the values of k and a_0H_0 as well, gives

$$N_* = 57.4 + \frac{1}{3}\ln\left(\frac{V_{\rm end}^{1/4}}{10^{16}\rm GeV}\right). \tag{13}$$

[1] A similar computation can be performed also in the case of more fermionic degrees of freedom. However, since here we are not discussing the details of the fermion sector phenomenology, but just its contribution to the effective potential, we limit ourselves to the minimal setup.

[2] There is also a RGE for y to be solved. In a minimal setup in which the Weyl fermion is only coupled to ϕ, the beta function for such a coupling would behave as $\beta_y \approx y^3$. If $y \ll 1$, then the running of y becomes negligible and y can be safely treated as a constant.

[3] The choice is just a convenient parameterisation. Even if we would assume $\lambda(M) \neq 0$, we can always find a new scale $M^* = M\exp(\frac{\lambda(M)}{2\beta})$ at which $\lambda(M^*) = 0$. Therefore the computations would then proceed in the same way from Eq. (5) with simply M^* in place of M.

[4] Indeed, we find $y = 4 \times 10^{-4}$ (see conclusions), which is not that far from the expected range.