Warm inflation in the light of swampland criteria

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Warm inflation seems to be the most befitting single-field slow-roll inflation scenario in the context of the recently proposed swampland criteria. We investigate the constraints these swampland criteria impose on warm inflation parameters and show that warm inflation is in accordance with both the current cosmological observations and the proposed swampland criteria in both weak and strong dissipative regimes depending on the value of the parameter c, which limits the slope of the inflaton potential according to the criteria.

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I. INTRODUCTION

At present, cosmic inflation [1–4] is dominantly believed to be the mechanism which provides the initial seeds of inhomogeneities for the observed cosmic microwave background anisotropies and the large scale structures. According to the basic picture of cosmic inflation, this brief period of accelerated (quasi)exponential (de Sitter) expansion of the Universe is driven by one slowly rolling scalar field, the inflaton. Such exponential expansion leaves the Universe devoid of any matter at the end of inflation, which calls for a period of reheating before the Universe reenters the standard big bang (decelerating) phase. Based on when and how the Universe is "reheated," the basic inflationary mechanism bifurcates: the original picture in which the inflaton field oscillates at the bottom of its potential at the end of inflation reheating the Universe by dissipating its energy to a radiation bath (we will refer to this mechanism as "cold inflation") and an alternative scenario in which the inflaton field keeps dissipating its energy to a radiation bath during the course of inflation maintaining a constant radiation energy and thus avoids the conventional "reheating phase" at the end of inflation (this mechanism is known as warm inflation (WI) proposed in Ref. [5]).

Inflation, believed to have taken place at the grand unified theory scale or below, is assumed to be described by low-energy effective field theories (EFTs). Such EFTs can be ultraviolet complete if they can be successfully embedded in a quantum theory of gravity, such as string theory. String theory provides large landscapes in which EFTs with Minkowski and anti-de Sitter vacua can be formulated with a consistent quantum theory of gravity, whereas EFTs with de Sitter vacua lie in the surrounding "swamplands" where EFTs coupled to gravity render quantum theory of gravity inconsistent. This has led to a set of criteria, like the weak gravity conjecture [6] and the recently proposed two swampland criteria [7], to ensure any (meta)stable de Sitter vacuum EFT does not lie in the desired string landscapes. These two swampland criteria, barring de Sitter vacuum from string landscapes, pose potential threats to the basic mechanism of slow-roll inflation (as has been observed in Ref. [8]), which we explain below.

We state the two swampland criteria as proposed in Ref. [7] and the reasons for concern raised by these criteria as far as inflationary dynamics is concerned:

(i) Swampland criterion I ($\mathcal{SC}$I).—This criterion puts an upper bound on the field range traversed by scalar fields in low-energy effective field theories as

$$\frac{\Delta\phi}{M_{\rm Pl}} < \Delta, \tag{1}$$

where $\Delta \sim \mathcal{O}(1)$ and $M_{\rm Pl}$ is the reduced Planck mass. This criterion emerges from the belief that there is a finite radius in field space of the EFT in which the effective Lagrangian remains valid. At large distances D, generation of a tower of light scalar modes with masses

$$m \sim M_{\rm Pl}\exp(-\alpha D), \tag{2}$$

with $\alpha \sim \mathcal{O}(1)$, renders the validation of the effective Lagrangian [8].

Lyth, in his seminal paper [9], devised a lower bound on the range traversed by the inflaton field

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over the course of (single-field) inflation, dubbed the Lyth bound, which is related to the ratio (r) of amplitudes of the tensor and the scalar perturbations produced during inflation and is stated as [10]

$$\frac{\Delta\phi}{M_{\rm Pl}} \gtrsim \Delta N\sqrt{\frac{r}{8}}, \tag{3}$$

where $\Delta N(\sim 60)$ is the number of e-folds of the duration of inflation. One should note that to obtain sub-Planckian field excursions during the course of inflation ($\Delta\phi \lesssim M_{\rm Pl}$), as demanded by the swampland criterion given in Eq. (1), one must have $r \lesssim \mathcal{O}(10^{-3})$ for $\Delta N \sim 60$ [11]. The recent observations by Planck and BICEP2/KEK put an upper bound on the tensor-to-scalar ratio as $r < 0.064$ [12].

Generic polynomial scalar potentials, like quartic and quadratic which appear in chaotic inflation models [3], are known to be yielding to large tensor-to-scalar ratios $[\mathcal{O}(10^{-1})]$ and hence are disfavored by the current data [12]. However, the "plateau models," like Higgs inflation [13], R^2 (Starobinsky) inflation [2], pole inflation [14], and α-attractor models [15], are known to yield such low tensor-to-scalar ratios ensuring small field excursions during the course of inflation and thus are not in much tension with $\mathcal{SC}$I.

(ii) Swampland criterion II ($\mathcal{SC}$II).—The second swampland criterion puts a lower bound on the gradient of the scalar field potentials of any EFT as

$$M_{\rm Pl}\frac{|V'|}{V} \gtrsim c, \tag{4}$$

where $c \sim \mathcal{O}(1)$. Here, the prime denotes a derivative with respect to the inflaton field. It is shown in Ref. [7] that the actual value of c depends on the details of compactification, and it often turns out to be of the order of $\sqrt{2}$ or greater in many string realizations and is not less than unity [16]. However, it has been argued in Ref. [17] that c as small as $\mathcal{O}(10^{-1})$ does not go against perceiving de Sitter vacua in String landscapes.

Single-field slow-roll inflationary dynamics, with a canonical kinetic term and Bunch-Davies vacuum state, falls short in three different ways of meeting this second swampland criterion:

(1) First of all, the slow rolling of the inflaton field is ensured by the flatness of its potential demanding the slow-roll parameters to be much smaller than unity. Thus, the slow-roll condition,

$$\epsilon_\phi \equiv \frac{M_{\rm Pl}^2}{2}\left(\frac{V'}{V}\right)^2 \ll 1, \tag{5}$$

itself is in direct conflict with $\mathcal{SC}$II.

(2) Second, single-field slow-roll inflation with a canonical kinetic term and Bunch-Davies initial states predicts

$$r = 16\epsilon_\phi. \tag{6}$$

Thus, $\mathcal{SC}$II demands that $r > 8c^2$, which is in conflict with the current observational upper bound on r as $r < 0.064$ [12] even if we consider c to be as low as 10^{-1} [17]. This rules out the standard chaotic inflation potentials[1] as well as the "plateaulike" potentials [2,13–15], which were in accordance with the $\mathcal{SC}$I as discussed above. This problem can be avoided in cases of non–Bunch-Davies initial states [19,20] and noncanonical kinetic terms such as in k inflation [21] models, which, due to the modified dynamics, introduces a suppressing factor on the rhs of the above equation, alleviating the constraint on r.

(3) Last, we note that, irrespective of the form of the potential during inflation, in the standard cold inflation scenario the inflaton field decays during reheating into the radiation bath while oscillating at the bottom of it potential, where $V' \sim 0$. This is also in conflict with the second swampland criterion [8].

It has been recently noted in Ref. [21] that warm inflation [5] suits $\mathcal{SC}$II the most. First of all, we note that as in WI the inflaton field dissipates to the radiation bath during the course of inflation, and such dynamics brings in an additional friction term to the inflaton slow-roll equation of motion,

$$3H\dot{\phi} + \Gamma\dot{\phi} \sim -V', \tag{7}$$

where Γ is the decay rate of the inflaton field. This yields the slow-roll condition as

$$\epsilon_\phi \ll 1 + Q, \tag{8}$$

where $Q = \Gamma/3H$. It is easily observed that in the strong dissipative regime ($Q > 1$) $\mathcal{SC}$II can be easily met without hampering the slow-roll condition of the inflaton field. Second, we note that, due to the modified inflaton dynamics, WI predicts [22–31]

$$r = \left(\frac{H}{T}\right)\frac{16\epsilon_\phi}{(1+Q)^{\frac{5}{2}}}, \tag{9}$$

where T is the temperature of the radiation bath with $T > H$. Thus, in the strong dissipative regime ($Q > 1$), the suppressing factor $(H/T)/(1+Q)^{5/2}$ helps evade $\mathcal{SC}$II to remain in tune with the present observations.

[1]It has been recently argued in Ref. [18] that chaotic inflation on a brane can be realized with polynomial potentials $V(\phi) \sim \phi^p$ for fine-tuned values of p as $p \lesssim 0.35$.

Last, we note that as WI, by construction, does not call for a reheating phase at the end of inflation, it also does not call for potentials with $V' \sim 0$, which further helps WI remain in accordance with $\mathcal{SC}$II.

The aim of the present article is to further investigate the WI scenario in the light of the swampland criteria to see how much these criteria constrain the parameters of WI, especially the parameter Q, which determines whether WI takes place in the weak $(Q < 1)$ or strong $(Q > 1)$ dissipative regime.

However, before approaching to the main analysis of the paper, it is important to analyze the warm inflationary dynamics in some detail as the basic mechanism of WI differs from that of conventional cold inflation to some extent. The WI scenario demands coupling of the inflaton field with light degrees of freedom (DOF) to which the inflaton field would dissipate its energy to maintain a constant radiation bath throughout the course of inflation. The presence of such light DOFs can potentially modify the inflaton potential which drives inflation. As the swampland criteria are all about bounds on the form of the scalar (inflaton) potentials in any EFT, it is important to scrutinize how much the inflaton potential gets modified due to the presence of such light DOFs in a warm inflationary model. It was realized soon after the proposal of warm inflation that coupling the inflaton with such light DOFs during inflation is indeed a taxing job [32], and the reason for it is twofold which can be easily understood if we consider Yukawa-like couplings $g\phi\bar{\psi}\psi$ of the inflaton field ϕ with light spinor fields ψ's. First of all, the inflaton induces large masses to the sprinor fields $m_\psi = g\phi$, and in return, the spinor fields induce large thermal corrections to the inflaton mass $(m_\phi \sim gT)$ barring slow roll for $T > H$. Both these hurdles can be overcome by fine tuning the coupling g, which then goes against the effective dissipation of the inflaton field to the light DOFs, and hence the scenario fails to sustain a constant thermal bath. This road block can be avoided in two circumstances. The first way out would be to allow the inflaton to couple to intermediate heavy scalars which will then eventually decay to the light DOFs. It was extensively estimated in Ref. [33] that in such warm inflationary scenarios the thermal correction to the inflaton potential is quite negligible within a global supersymmetry setup. The WI with brane construction setup has also been studied in Ref. [34] in this context, and it has been shown that the thermal corrections turn out to be Boltzmann suppressed. The other scenario in which warm inflation can be successfully realized is the "warm little inflaton" scenario [26,31], in which the inflaton is treated as pseudo–Nambu-Goldstone boson [corresponding to the relative phase between two complex Higgs scalars that collectively break a local $U(1)$ symmetry] coupled to a pair of fermionic fields through Yukawa interactions. The advantage of this scenario is that it bounds the masses of the fermions as $gM\cos(\phi/M)\bar{\psi}_1\psi_1$ and $gM\sin(\phi/M)\bar{\psi}_2\psi_2$ (with M as the vacuum expectation value of the two Higgs scalar) and also, for $m_{\psi_{1,2}} \ll T$, the thermal mass correction due to the fermions cancels between the contributions of both the fermions, leaving only the subleading Coleman-Weinberg term. To achieve such cancellations, discrete exchange symmetry is imposed in the scalar-spinor sector. The zero-temperature form of the potential is also protected against large thermal correction due to the gauge symmetry of the underlying theory. Hence, in this case as well, the thermal corrections modify the slope of the inflaton potential negligibly and thus do not affect the slow-roll conditions of the warm inflation scenario (see Ref. [35], for example), leaving the swampland constraints intact for the warm inflationary case.

II. SWAMPLAND CRITERIA AND WARM INFLATION

We will first analyze $\mathcal{SC}$I to see what constraints it imposes on the WI parameter Q. As in cold inflation, the $\mathcal{SC}$I is in contrast with the Lyth bound, as stated in Eq. (3), and we need to determine the Lyth bound in the context of WI to appraise $\mathcal{SC}$I. We first note that the scalar power spectrum in WI receives two additive factors along with the form we get in cold inflation as [25]

$$P_{\mathcal{R}} = \left(\frac{H}{\dot{\phi}}\right)^2 \left(\frac{H}{2\pi}\right)^2 \left[1 + 2n + \left(\frac{T}{H}\right)\frac{2\sqrt{3}\pi Q}{\sqrt{3+4\pi Q}}\right], \qquad (10)$$

where the last factor appears due to the presence of dissipative term in the inflaton equation of motion and $n \equiv 1/(\exp(-k/aT) - 1)$ denotes the Bose-Einstein distribution of the thermalized inflaton fluctuations. However, when $T > H$, a condition for thermal equilibrium, the scalar power spectrum can be approximated as [24]

$$P_{\mathcal{R}} \approx \frac{1}{8\pi^2\epsilon_\phi}\frac{H^2}{M_{\rm Pl}^2}(1+Q)^{5/2}\left(\frac{T}{H}\right). \qquad (11)$$

The weakly coupled tensor modes, however, remain unaffected by the WI dissipative terms, yielding the same tensor spectrum as in cold inflation:

$$P_T = \frac{2H^2}{\pi^2 M_{\rm Pl}^2}. \qquad (12)$$

These yield the tensor-to-scalar ratio in the warm inflation scenario as

$$r = \left(\frac{H}{T}\right)\frac{16\epsilon_\phi}{(1+Q)^{\frac{5}{2}}}, \qquad (13)$$

which can also be written as

$$r = \left(\frac{H}{T}\right)\frac{16\epsilon_H}{(1+Q)^{\frac{3}{2}}}, \tag{14}$$

as in WI,

$$\epsilon_H \equiv -\frac{\dot{H}}{H^2} \approx \frac{\epsilon_\phi}{1+Q}. \tag{15}$$

The modified inflaton dynamics of WI suggests

$$\dot{\phi} = -\frac{\sqrt{2\epsilon_H}}{\sqrt{1+Q}} M_{\mathrm{Pl}} H, \tag{16}$$

which yields

$$\frac{\Delta\phi}{M_{\mathrm{Pl}}} = \sqrt{\frac{r}{8}\left(\frac{T}{H}\right)}(1+Q)^{\frac{1}{2}}\Delta N, \tag{17}$$

rendering the modified Lyth bound for warm inflation as [36,37]

$$\frac{\Delta\phi}{M_{\mathrm{Pl}}} \gtrsim \sqrt{\frac{r}{8}\left(\frac{T}{H}\right)}(1+Q)^{\frac{1}{2}}\Delta N. \tag{18}$$

$\mathcal{SC}$I suggests scalar field excursions (in Planck units) to be smaller than $\Delta \sim \mathcal{O}(1)$. But as the actual value of the parameter Δ is uncertain, we simply demand sub-Planckian field excursions during the 60 e-folds duration of inflation. Thus, demanding $\Delta\phi/M_{\mathrm{Pl}} < 1$, we get

$$r < \frac{8}{(\Delta N)^2\sqrt{1+Q}}\left(\frac{H}{T}\right). \tag{19}$$

But the current observational bound suggests that $r < 0.064$. These two conditions can be simultaneously met if

$$\frac{1}{(\Delta N)^2\sqrt{1+Q}}\left(\frac{H}{T}\right) < 0.008, \tag{20}$$

yielding

$$1+Q > \frac{1}{(0.008)^2(\Delta N)^4}\left(\frac{H}{T}\right)^2 \sim 10^{-5}, \tag{21}$$

where we have considered $\Delta N \sim 60$ and $H/T \sim 10^{-1}$. We note that $\mathcal{SC}$I is easily satisfied by WI even in the weak dissipative regime $Q < 1$.

Let us now consider the second swampland criterion, which calls for the following condition:

$$\epsilon_\phi > \frac{c^2}{2}. \tag{22}$$

Combining this condition with Eq. (13), we get

$$r > \left(\frac{H}{T}\right)\frac{8c^2}{(1+Q)^{\frac{5}{2}}}. \tag{23}$$

This condition satisfies the current observational upper bound on r if

$$\left(\frac{H}{T}\right)\frac{c^2}{(1+Q)^{\frac{5}{2}}} < 0.008, \tag{24}$$

yielding

$$1+Q > \left(\frac{H}{T}\right)^{\frac{2}{5}}\left(\frac{c^2}{0.008}\right)^{\frac{2}{5}} \sim 4, \tag{25}$$

where we have considered $c \sim \sqrt{2}$ and $H/T \sim 10^{-1}$. We note that $\mathcal{SC}$II puts a more stringent bound on Q than $\mathcal{SC}$I. $\mathcal{SC}$II demands WI takes place in the strong dissipative regime $Q > 1$. The upcoming observations, like COrE [38] and LiteBIRD [39], will search for tensor-to-scalar ratio $r \sim \mathcal{O}(10^{-3})$, and a nonobservance of r by such missions would drive WI further deep into the strong dissipative regime. On the other hand, if c turns out to be of the order of 10^{-1}, as has been argued in Ref. [17], then we can see that with the present bound on r one gets $1+Q > 0.5$. In such a situation, even weak dissipative regime WI scenarios would be in accordance with both the swampland criteria as well as with the current observations.

III. DISCUSSION AND CONCLUSION

It was readily realized after the recently proposed swampland criteria [7] that these criteria regarding the formulation of UV-complete low-energy EFTs might have severe consequences for the cosmological epochs relying on the de Sitter vacuum, like inflation and dark energy [8]. Since then, a series of papers has been written to counter such stumbling blocks and to put both inflation and dark energy back on track; a nonexhaustive list of such analysis would include Refs. [16–21,40–42]. Reference [21] appraises several single-field slow-roll scenarios in the light of the second swampland criterion and shows that warm inflation, among all, turns out to be the best inflationary scenario befitting this criterion.

In this paper, we analyzed both swampland criteria in the context of WI to investigate what constraints the criteria can impose on WI parameters, in particular Q, which serves as an indicator of whether WI takes place in the weak ($Q < 1$) or strong ($Q > 1$) dissipative regime. We found that $\mathcal{SC}$II puts a more stringent bound on Q than $\mathcal{SC}$I, implying that WI should take place in the strong dissipative regime with $Q > 3$ when $c \geq \sqrt{2}$. However, if the parameter c, the actual value of which depends on the methods of compactifications, can be brought down to as low as 10^{-1}, then $\mathcal{SC}$II can also allow for weak dissipative regimes in WI as far as the present observations are concerned. We also note that if future observations lower the upper bound on the tensor-to-scalar ratio then $\mathcal{SC}$II would push WI deeper into the strong dissipative regime.

While this paper was under preparation, a similar analysis was presented in Ref. [43]. One should note that the conclusion drawn in this paper differs from that of Ref. [43] as the latter concludes that Q should be of the order $\Delta N \sim 60$ or larger for WI to evade the swampland criteria, driving WI models very deep into the strong dissipative regime. This is of concern, as it was shown in Ref. [43] that most of the models of WI do not allow for such large Q as that would result in a redder scalar spectrum, which is not in accordance with the observation of the scalar spectral tilt n_s. Our conclusion differs from that of Ref. [43], as we showed that the required value of Q to be in tune with the swampland criteria is an order of magnitude smaller than $\Delta N \sim 60$ and the swampland criteria can even allow for the weak dissipative regime in WI if the parameter c appearing in $\mathcal{SC}$II turns out to be smaller than unity (but positive). Hence, we note that WI still remains the best possibility among the single-field slow-roll inflation scenarios if the swampland criteria stand the test of time.

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