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Inflation at the GUT scale in a Higgsless universe

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We revisit inflation in induced gravity. Our focus is on models where the low scale Planck mass is completely determined by the breaking of the scaling symmetry in the field theory sector. The Higgs-like field which breaks the symmetry with a GUT-scale vacuum expectation value (vev) has nonminimal couplings to the curvature, induced by the gravitational couplings of the other light fields in the theory, so that its vev controls the gravitational strength. This field can drive inflation and give a low energy universe in very good agreement with the cosmological observations. The low energy dynamics of the standard model cannot be unitarized by the Higgsflaton, which decouples from the low energy theory, both because it picks up a large mass and because its direct couplings to the low energy modes are weakened. Instead, the short distance behavior of the standard model may be regulated by the dynamics of other light degrees of freedom, such as in Higgsless models.

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The standard model (SM) of particle physics has been a singularly successful framework for explaining the observed dynamics of elementary particles. Its low energy canonical spectrum contains[1] 90 fermionic degrees of freedom, 27 vector bosons, with or without mass terms, and a single scalar degree of freedom which has so far eluded all attempts at detection. This evasive mode—the Higgs field—is special in many respects. It is the only fundamental scalar in the SM, having so far completely avoided direct detection. On the other hand, the whole structure of the SM hinges on its existence, because it is responsible for the unitarization of the electroweak sector of the theory and the generation of vector boson and fermion masses. Indeed, the Higgs spontaneously breaks the electroweak gauge symmetry, setting the mass scale of the massive SM fields.[2] The residual Higgs fluctuations then regulate the massive low energy electroweak sector and unitarize its scattering amplitudes. However, not all is well in the SM tale. As is well known, the Higgs mass is not radiatively stable, and its phenomenological value of $\sim$100 GeV, and consequently a flat potential, begs the question about what may possibly keep it there. This single missing SM degree of freedom is so vital to the whole model that a spectacular machine such as the Large Hadron Collider (LHC) has as one of its key tasks seeking for it, and for the physics which makes it possible.

But, what if the LHC does not find a fundamental scalar? An absence of the Higgs would push us into alternative explanations of the observed low energy SM dynamics. Models without a light Higgs may have to separate the origin of masses from the new physics which unitarizes the low energy theory. Indeed, what would (not) be directly observed at the LHC are only the fluctuations of the Higgs field, and not its zero mode. Examples where the unitarization of SM amplitudes is disentangled from the origin of mass have already been considered, and among them recently the Higgsless models in extra dimensions [1] attracted much attention. These models are realized as brane setups in cutoff AdS space [2], with AdS radius L, with fermions localized on branes and gauge bosons propagating in the bulk. In the effective $4D$ theory this yields towers of Kaluza-Klein gauge bosons, whose masses and couplings are determined by the warping of the bulk AdS geometry. In the dual cutoff AdS/CFT, they are duals of light conformal field theory (CFT) states, with masses below the CFT UV cutoff $\mu \sim 1/L$, whose number is $\mathcal{N} \sim (M_4 L)^2$ [3]. Curiously, this counting of light states agrees with the recent ideas on a relation of the number of light states below some UV cutoff and hierarchy between this cutoff and the Planck scale [4]. In fact, such models are naturally related to technicolor models, but now in the strong coupling regime as defined by way of the AdS/CFT duality, where the SM is also unitarized without invoking fundamental scalars.

Thus if no fundamental scalar were observed at the LHC, we will have had the (poor) consolation of verifying experimentally the theoretical prejudice against light fundamental scalars. Such an outcome would speak loudly against the existence of *any* light fundamental scalar, indicating that nature may choose other routes for realizing the low energy SM. We stress once again here, that while

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[1]The discovery of the neutrino masses has already taken us outside of the canonical SM, extending the spectrum by at least as many as 6 more fermions.

[2]With the possible exception of neutrinos, which by the seesaw mechanism may inherit masses from dynamics at much higher scales.

this may be an extreme point of view, it is not *yet* excluded by any experimental facts. SM with the Higgs is the simplest means of describing the observed particle dynamics, but the Higgs is still missing, and its existence has been questioned. Thus it is of interest to consider other implications of a missing Higgs, particularly as it relates to the paradigm of naturalness. While speculative from the point of view of our current expectations, they yet remain to be excluded.

Indeed, the absence of the Higgs would have other implications if we take the view that its presence were to support naturalness. Beyond the SM physics, scalars also play a key role in cosmology, where they are prototypical inflatons. A scalar field provides the simplest dynamics necessary to inflate the universe, ensuring that its large smooth and nearly flat swaths survive to the present epoch [5]. The scalars can be inflatons if their potential is very flat in the units of the natural cutoff, compared to which they are light. This is qualitatively similar to the SM Higgs, although in practice the requirement for potential flatness is quantitatively weaker. Nonetheless, all inflaton models need some amount of fine-tuning to make sure the potential remains flat in spite of the couplings to other matter, necessary for reheating.

In the event that LHC finally discovers the Higgs, it will be easy to imagine that other light scalars with flat potentials appear in nature, regardless of why that may be so. One could be the inflaton, at a scale well below the cutoff, and well separated from UV physics, and ultimately gravity. On the other hand, if no light scalars are seen, a logical consequence may be that light scalars are hard to sustain. In such an instance, the scalars would drift up to near the cutoff, which may be at the grand unified theory (GUT) scale $\sim 2 \times 10^{16}$ GeV, as hinted at already from sub-TeV-scale physics explored so far, by the proton stability, the seesaw explanation neutrino masses, and the clues from gauge coupling unification.[3] So if the scalars cannot be stabilized near the TeV scale, radiative stability may be attained if the scalar masses are pushed high, to near the GUT scale, where even the mode responsible for inducing SM masses may end up. In such a scenario, the inflaton would be no exception. However unlikely this option may now seem, the conspicuous absence of the Higgs from the observed bestiary of elementary particles found to date points to the fact this is not yet *impossible*. Indeed, the Higgsless models of various kinds already account for this in the SM sector. In this note, we shall outline how to allow for inflation in such a universe, basing it on a Higgs-like field, which spontaneously breaks the scale symmetry at the GUT scale, and gets a flat enough effective potential, as it induces the (small) hierarchy between the GUT scale and the Planck scale.

In the early days of inflationary model building, the possibility of driving inflation by the SM Higgs has been tried, but without immediate success [7]. With minimal couplings to gravity, the SM constraints force the scalar self-couplings to be too large to yield satisfactory inflationary density perturbations (see [8] for a review). These problems can be ameliorated if nonminimal couplings to gravity are allowed. In particular, in the induced gravity framework [9,10] one can get the right density perturbations even if the scalar self-couplings are much larger than in the minimal coupling case [11]. Recently it has been noted [12] that if the scalar has direct coupling to the curvature, $\sim \xi |\phi|^2 R$, and there is also the standard Einstein-Hilbert term in the theory, $\sim M_{\rm Pl}^2 R$, then the scalar could both drive a low scale inflation, yielding the right density contrast, and serve as the Higgs after inflation. For this claim, it is crucial that the gravitational sector contains the Einstein-Hilbert term.[4] If it were not so, the cosmic background explorer (COBE) normalization and the phenomenologically required Higgs vev, $\langle H \rangle \sim 246$ GeV, would force the value of the Planck scale to be at the ~ 10 TeV scale.

Our route here is very different. We imagine that the underlying theory is conformal in the UV, including the gravitational sector. This means that the bare gravitational Einstein-Hilbert term is absent from the action, which instead contains higher derivative terms, starting with the curvature squared invariants,

$$\mathcal{S} = \int d^4x \sqrt{g} \Big[A\mathcal{GB} + BC^2_{\mu\nu\lambda\sigma} + CR^2 - \frac{1}{2}(\nabla\phi)^2 - \frac{\lambda}{4}(\phi^2 - v^2)^2 - \mathcal{L}_{\rm matter}(g^{\mu\nu}, \phi, \psi) + \ldots \Big], \tag{1}$$

where $\mathcal{GB} = R^2_{\mu\nu\lambda\sigma} - 4R^2_{\mu\nu} + R^2$ is the Gauss-Bonnet combination, $C_{\mu\nu\lambda\sigma}$ is the Weyl tensor, and A, B, C are some dimensionless constants. This theory is in fact renormalizable, as shown some time ago in [14] and, later, in works on induced gravity [10] and relation between Newton's constant and scale symmetry breaking [15]. On the other hand, suppose there is a somewhat large number of degrees of freedom in the matter sector, $\sim \mathcal{O}(10^4)$, including those which will become the low energy SM. If there is a gauge group in the theory which confines at some scale, dimensional transmutation will yield an IR cutoff, which will be fed back to the scalars. There may also be explicit symmetry breaking terms in the scalar sector, with the scalars which are not protected from radiative corrections from the strong gauge group.

Either way, the matter sector quantum field theory will be characterized by a dimensional cutoff. Then, the quantum one-loop effects will generate contributions to the

[3] Some features of dynamics in Higgsless models as pertaining to these scales, and specifically issues of relevance for unification have been addressed in [6].

[4] More aspects of this scenario were considered in [13].

action of the form $\sim\Lambda^2 R$ [9,10]. In general, these corrections will depend on the cutoff itself, as well as the value of field vevs around which the corrections are calculated. We will *assume* that the field independent contributions to Λ can be neglected. This could be justified as follows. The quantum contributions to $\sim R$ term will come as $(\Lambda^2 + c\phi^2)R$ from every degree of freedom which couples to gravity. If these degrees of freedom are all weakly coupled, one would expect that the bare cutoff terms may dominate. On the other hand, if some are in strong coupling, the strong coupling effects may conspire between different orders in the loop expansion and retain the appearance of conformality, such that the dimensional transmutation which they trigger may occur at a scale well below the strong coupling scale [16]. Thus this scale could be smaller than the one directly sampled by the Higgs symmetry breaking.[5] Then the leading order contributions to Λ may come from the IR masses of the fields residing in the geometry, yielding by linear superposition $\Lambda^2 \sim \sum_k m_k^2$. If these masses are generated directly by a symmetry breaking induced by a Higgs-like field (Higgs for short from now on), $m_k(\phi) \sim g\phi$, this would yield $\Lambda^2 \simeq \mathcal{N} g^2\phi^2$, which can dominate over the hard cutoff contributions. Here for simplicity we assume that all the Yukawa couplings are approximately the same. The number $\mathcal{N}$ counts the fields in the theory which are Higgsed by ϕ, and so this yields $\xi \sim \mathcal{N} g^2$. Again, this is consistent with the recent ideas about the large number of light fields inducing the hierarchy between the mass scale where they reside and the Planck scale [4], although it would be a much more conservative quantitative implementation of such a framework. Note that the crucial aspect of this idea is that the conformal symmetry breaking which induces the Einstein-Hilbert term is *soft*, in that the hard cutoff contributions must be subleading, which typically may not occur in weak coupling [16].

Of course, the scalar which breaks the symmetry cannot be the usual Higgs [12,17], since its mass will be too large, as would be natural by the low energy accounting of radiative corrections. This scalar will have its mass and vev set by the scale where the conformal symmetry breaks down. To reflect this, we will dub it the "Higgsflaton," and take the symmetry breaking scale to be the GUT scale.[6] However the crucial property that allows the Higgsflaton to drive inflation, and therefore get a somewhat flatter potential, is its coupling to the Ricci scalar. The key reason is that the parameter ξ, of the order of 10^4, needed to induce the hierarchy between the GUT scale and the Planck scale, also seesaws the scalar mass from the GUT scale down to $m_\varphi \sim v/\sqrt{\xi}$, flattening the scalar potential just enough. Moreover, this number precisely reproduces the COBE normalization[7] of the scalar density perturbations in this model. Given the argument for how the Einstein-Hilbert term comes about, the value of ξ can be obtained by positing that the Higgsflaton gives mass to about 10^4 degrees of freedom, with Yukawa couplings $g \sim 1/3$, which therefore live at the GUT scale, and whose loops induce the Einstein-Hilbert term. In this case, the low energy standard model is unitarized by some other degrees of freedom, e.g. as in the Higgsless models [1]. Note that in this scenario—as in the Higgsless model—we are not addressing the origin of the electroweak scale, which should be attributed to some other strong dynamics that does not necessarily involve scalar modes. At least the SM fields, being outnumbered by the other degrees of freedom in the theory, and much lighter than most, will not contribute significantly to the generation of the Einstein-Hilbert term, which would be largely insensitive to their presence.

Let us now outline the cosmological scenario. In light of the discussion above, the low energy theory, below the scale symmetry breaking, is given by the effective $4D$ action

$$\mathcal{S} = \int d^4x\sqrt{g}\Big[\frac{1}{2}\xi\phi^2 R - \frac{1}{2}(\nabla\phi)^2 - \frac{\lambda}{4}(\phi^2 - v^2)^2 - \mathcal{L}_{\text{matter}}(g^{\mu\nu}, \phi, \psi) + \ldots\Big], \tag{2}$$

where $\mathcal{L}_{\text{matter}}$ includes the standard model and additional matter fields which unitarize it at the $\sim$TeV scale, collectively denoted by ψ, and ϕ is the Higgsflaton scalar field modulus, with a nonminimal coupling to curvature $\xi\phi^2 R$. The ellipsis stand for additional terms which we assume to be mostly negligible. The Higgsflaton phase is in $\mathcal{L}_{\text{matter}}$ as a longitudinal component of a gauge boson, so that $\mathcal{L}_{\text{matter}}$ is written in a unitary gauge. With the assumptions above, the parametrization of its leading order low energy dynamics by (2) is accurate in the limit of weak gravity. On the other hand, although in the regime of background field values $\phi \sim 0$ the field theory in (2) is perturbative, gravity as encoded by (2) becomes strong. So sufficiently close to the origin in field space the theory cannot be described by (2). However, in this regime the scale symmetry is restored, and the gravitational theory reverts back to the curvature squared action, with a negligible Einstein-Hilbert correction.

At any rate, at low energies for large values of ϕ which break the symmetry, gravity will be weak when $\xi \gg 1$. In this limit, we can use the field equations derived from (2) to describe the background geometry. At a minimum,

[5] At least in the weak coupling this may occur, as we know from the example of QCD, where low energy quark masses are mainly attributed to the electroweak symmetry breaking.

[6] The proximity of the GUT scale to the Planck scale makes the presence of fundamental scalars near the GUT scale appear more plausible, since at those scales one may get away without mechanisms that protect their masses from radiative corrections.

[7] Possible connections between the GUT scale and primordial density perturbations were noted in [18], albeit realizations were different.

$\phi = \pm v$, if we integrate out the scalar the theory reduces to $\mathcal{S}_{\rm eff} = \int d^4x\sqrt{g}(\frac{1}{2}\xi v^2 R - \mathcal{L}^{\rm eff}_{\rm matter}(g^{\mu\nu}, \psi) + \ldots)$, which shows that the effective low energy Planck scale around the scalar vacuum is

$$M_{\rm Pl}^2 = \xi v^2. \tag{3}$$

To see the scalar dynamics we can go to the unitary gauge where all fields are canonically normalized. Taking the conformal transformation and scalar field redefinition [11,19,20],

$$\hat{g}_{\mu\nu} = \left(\frac{\phi}{v}\right)^2 g_{\mu\nu}, \qquad \varphi = M_{\rm Pl}\sqrt{6 + \frac{1}{\xi}}\ln\left(\frac{\phi}{\phi_0}\right), \tag{4}$$

where ϕ_0 is an arbitrarily chosen normalization, yields the Einstein frame action

$$\begin{aligned}\mathcal{S} = \int d^4x\sqrt{\hat{g}}\Big\{&\frac{M_{\rm Pl}^2}{2}\hat{R} - \frac{1}{2}(\hat{\nabla}\varphi)^2 - \hat{V}(\varphi)\\ &- \left(\frac{v}{\phi_0}\right)^4 e^{-4(\varphi/M_{\rm Pl}\sqrt{6+1/\xi})}\\ &\times \mathcal{L}_{\rm matter}((\phi_0/v)^2 e^{2(\varphi/M_{\rm Pl}\sqrt{6+1/\xi})}\hat{g}^{\mu\nu}, \varphi, \psi) + \ldots\Big\}.\end{aligned} \tag{5}$$

The new effective potential is, using Eq. (3),

$$\begin{aligned}\hat{V}(\varphi) &= \frac{\lambda}{4}\frac{(\phi^2 - v^2)^2}{(\phi/v)^4}\\ &= \frac{\lambda M_{\rm Pl}^4}{4\xi^2}\left[1 - \left(\frac{v}{\phi_0}\right)^2 e^{-2(\varphi/M_{\rm Pl}\sqrt{6+1/\xi})}\right]^2.\end{aligned} \tag{6}$$

The minima $\phi = \pm v$ clearly correspond to $\varphi = M_{\rm Pl}\sqrt{6 + \frac{1}{\xi}}\ln(\frac{v}{\phi_0})$. Around the minimum, the curvature of the effective potential (6) yields the scalar mass

$$m_\varphi^2 = \partial_\varphi^2 \hat{V} = \frac{2\lambda M_{\rm Pl}^2}{\xi^2(6 + 1/\xi)}, \tag{7}$$

that, together with Eq. (3), implies

$$m_\varphi^2 \simeq \frac{\lambda v^2}{3\xi}, \tag{8}$$

in the limit when $\xi \gg 1$. Obviously, in the limit $\xi \sim 1$, $m_\varphi \sim M_{\rm Pl}$ and so this case is less interesting. This is precisely the seesaw effect in the scalar sector, which we alluded to in the introductory discussion. Indeed, that this is akin to seesaw can be seen by eliminating ξ from Eq. (8) by using Eq. (3), which yields

$$m_\varphi^2 \simeq \frac{\lambda v^4}{3M_{\rm Pl}^2}, \tag{9}$$

precisely a seesaw mass formula. In fact, the dynamics responsible for flattening the potential is conceptually similar to scalar "seizing" of [21], except that the large wave function renormalization involves the graviton as well as the scalar field.

We note that the "strong gravity regime" $\phi \sim 0$ in the Einstein frame variables corresponds to the limit $\varphi \to -\infty$, where the potential (6), and also all mass scales in the matter sector in (5) diverge. This of course is simply the restatement of the fact that the ratio of any mass scale μ and the effective Planck mass $M_{\rm Pl} = \sqrt{\xi}\phi$ diverges when $\phi \to 0$. This manifestly excludes the limit $\varphi \to -\infty$ from the low energy action (5), because in this case one must restore the quadratic curvature terms which were ignored in writing the effective action (2).

For the potential (6), clearly inflation occurs when $|\phi| > v$. In this limit gravity is weak, and furthermore the potential behaves like a cosmological constant. This can be readily seen from (6), since when $|\phi| > v$, $\hat{V} \to \frac{\lambda v^4}{4} = \frac{\lambda M_{\rm Pl}^4}{4\xi^2}$. Thus, since the potential asymptotes a constant when $\varphi \to \infty$, which smoothly goes to the minimum $|\phi| = v$, sufficient inflation followed by a graceful exit will occur when φ is initially large. Note, however, that by the formula (9), the mass of the Higgsflaton at the minimum is comparable to the Hubble scale during inflation, so the slow roll may extend even as the field approaches the minimum. In the original variables, the initial value of the field ϕ need not exceed $M_{\rm Pl}$ when $\xi \gg 1$. This is qualitatively similar to assisted inflation [22], where the expectation value of the inflaton during inflation also need not be trans-Planckian. For more complicated potentials, which may even involve bigger powers of ϕ, however, the effective potential (6) will still typically have a maximum, and decay back to zero for very large values of φ. In such cases, it is still possible to have inflation if the initial value of φ will be near the maximum, which is expected to occur somewhere due to the random distribution of initial values [23].

Taking the background to be a spatially flat Friedmann-Robertson-Walker spacetime, we can use the slow-roll equations to describe the geometry at large scales. This yields

$$\begin{aligned}H^2 &\cong \frac{\lambda M_{\rm Pl}^2}{12\xi^2}\left[1 - \left(\frac{v}{\phi}\right)^2\right]^2,\\ \dot{\varphi} &\cong -\frac{2}{\sqrt{3}\sqrt{6 + 1/\xi}}\frac{\sqrt{\lambda}}{\xi}\left(\frac{v}{\phi}\right)^2 M_{\rm Pl}^2.\end{aligned} \tag{10}$$

Using these Eqs. (10), and recalling that curvature perturbations are independent of the conformal frame in which they are calculated [20,24], it is straightforward to determine the amplitudes of scalar and tensor perturbations generate during inflation. Their powers are, respectively,

$$\Delta_{\mathcal{R}}^2 = \left(\frac{H^2}{2\pi\dot{\varphi}}\right)^2 \cong \frac{\lambda}{128\pi^2\xi^2}\left(\frac{\phi}{v}\right)^4\left[1 - \left(\frac{v}{\phi}\right)^2\right]^4, \tag{11}$$

$$\Delta_h^2 = 8\left(\frac{H}{2\pi M_{\rm Pl}}\right)^2 \cong \frac{\lambda}{6\pi^2\xi^2}\left[1-\left(\frac{v}{\phi}\right)^2\right]^2, \qquad (12)$$

where we are taking the limit $\xi \gg 1$. Now, to determine the scale at which (11) and (12) need to match to the observed anisotropies, we need to relate the field values to the inflationary clock readings, conveniently given by the number of efolds before the end of inflation. Inflation will end when the field rolls near the minimum $|\phi| \simeq v$. However, to get a precise location of the end of inflation, we can use the slow-roll parameters in the Einstein frame, which for the potential (6) are

$$\epsilon = \frac{M_{\rm Pl}^2}{2}\left(\frac{\partial_\varphi \hat{V}}{\hat{V}}\right)^2 \simeq \frac{4}{3}\left[1-\left(\frac{v}{\phi}\right)^2\right]^{-2}\left(\frac{v}{\phi}\right)^4, \qquad (13)$$

$$\eta = M_{\rm Pl}^2 \frac{\partial_\varphi^2 \hat{V}}{\hat{V}} \simeq \frac{4}{3}\left[1-\left(\frac{v}{\phi}\right)^2\right]^{-2}\left[2\left(\frac{v}{\phi}\right)^2 - 1\right]\left(\frac{v}{\phi}\right)^2. \qquad (14)$$

Inflation will end when either ϵ or η become $\mathcal{O}(1)$. From (13), this will occur at $|\phi| \equiv \phi_* \simeq 1.47v$. This means that between some value $\phi > \phi_*$ and this terminal value ϕ_*, the universe will undergo N efolds of inflation, where N is related to ϕ according to

$$N = \int_{\phi_*}^{\phi} \frac{H}{\dot\varphi}\frac{d\varphi}{d\phi}d\phi \cong \frac{3}{4}\left(\frac{\phi}{v}\right)^2 - \frac{3}{4}\left(\frac{\phi_*}{v}\right)^2 - \frac{3}{2}\ln\left(\frac{\phi}{\phi_*}\right)$$
$$\simeq \frac{3}{4}\left(\frac{\phi}{v}\right)^2 - 1 - \frac{3}{2}\ln\left(\frac{\phi}{\phi_*}\right), \qquad (15)$$

where we have used (4), (13), and (14). Since the pivot scale where cosmic microwave backgrounds (CMB) observations are matched to the theory is $N_p = 55$, this implies that the formulas for amplitude of perturbations (11) and (12) read

$$\Delta_{\mathcal{R}}^2 \simeq \frac{\lambda}{72\pi^2\xi^2}(N+4.3)^2, \qquad \Delta_h^2 \simeq \frac{\lambda}{6\pi^2\xi^2}, \qquad (16)$$

for $N \simeq 55$. They yield $\Delta_{\mathcal{R}}^2 \simeq 4.9 \times \frac{\lambda}{\xi^2}$ and the tensor-to-scalar ratio $r = \Delta_h^2/\Delta_{\mathcal{R}}^2 \simeq 0.003$. This is within the reach of future observational confirmation by planned experiments of B-mode polarization observation of CMB such as B-Pol. Matching the curvature perturbations to the observed value of $\Delta_{\mathcal{R}}^2 = 2 \times 10^{-9}$ gives

$$\frac{\sqrt{\lambda}}{\xi} \simeq 2.0 \times 10^{-5}. \qquad (17)$$

We can also easily calculate the spectral index of the scalar perturbations. The standard formula $n_s = 1 + \frac{d\ln\Delta_{\mathcal{R}}^2}{d\ln k}$ gives

$$n_s \cong 1 - \frac{2}{N+4.3}, \qquad (18)$$

which translates numerically to $n_s = 0.97$, in excellent agreement with the CMB data. Aspects of the CMB constraints on the perturbations in the model based on (2) were also considered in [25].

What of particle physics scales in this theory? As it manifest from Eq. (17), inflationary dynamics constrains the ratio of the coupling constants λ and ξ. To break this degeneracy we can take the coupling λ to be perturbative, but not tiny, in order to relax the usual severe tunings in the field theory sector of the inflaton [11]. So, suppose that $\lambda \sim 10^{-2}$. In this case, Eq. (17) implies $\xi \sim 5000$, and so by Eq. (3) we find

$$v = \frac{M_{\rm Pl}}{\sqrt{\xi}} \sim 3 \times 10^{16}\ \text{GeV}, \qquad (19)$$

i.e. we find $v \simeq M_{\rm GUT}$, exactly as we asserted in the introductory discussion. In turn the Higgsflaton mass (7) in the vacuum $|\phi| = v$ is $m_\varphi = \sqrt{\lambda}v/\sqrt{3\xi} \sim 3 \times 10^{13}$ GeV, by Eq. (8), which thanks to the seesaw induced by the large parameter ξ is significantly below the symmetry breaking scale v.

As the field rolls down the slope of (6) towards the minimum, it passes through an inflection point and the local curvature of the potential, negative up on the plateau, will increase slowly, eventually ending inflation. After falling down the precipice to the potential well around the minimum, the field oscillates around it on a time scale of the order of m_φ^{-1}, reheating the universe. The details of reheating depend on the couplings of the Higgsflaton to matter. The simplest case is when in the original, Jordan frame, ϕ couples the SM fermions only via Yukawa couplings. In this case, the classical scaling symmetry allow us to completely remove and decouple the canonical Higgsflaton field φ from matter. To see this consider the transformation of the Jordan frame Yukawa term $\sqrt{g}\phi\bar\psi\psi$ under conformal transformation. The fermions will scale according to $\psi = (\phi/v)^{3/2}\Psi$, so that Ψ is canonically normalized, which turns Yukawa couplings into simple mass terms $\sqrt{\hat g}v\bar\Psi\Psi$ [26]. Without other direct couplings of φ to matter, reheating may occur in two stages. In the first stage, φ oscillates about the bottom of its potential and its self interactions rapidly lead to resonant amplification of the nonzero modes of φ, which rescatter on the surviving part of the condensate, eventually disrupting it [27,28], and ensuring that the universe is filled almost exclusively by quanta of φ with a typical momentum of $\mathcal{O}(m_\varphi)$. Subsequently, the quanta of φ will scatter against each other, in processes like $\varphi\varphi \to \Psi\Psi$ mediated by gravitons, and produce the SM matter. A typical time scale for this process is the gravitational scattering scale $\tau_{\rm gs} \simeq M_{\rm Pl}^4/m_\varphi^5$, which with the mass scale $m_\varphi \sim 3 \times 10^{13}$ GeV yields a reheating temperature $T_{\rm RH} \sim g_*^{-1/4}(M_{\rm Pl}/\tau_{\rm gs})^{1/2} \sim$ TeV.

In reality, however, the reheating will be more efficient, because there will be additional couplings. To start with, one-loop corrections will spoil the exact cancellation between the rescaling factors in Yukawa terms, yielding a leftover field-dependent mass $m_\Psi \sim v(\phi/v)^d$, where $d \sim \mathcal{O}(1)\frac{g^2}{4\pi^2}$ arises from the anomalous dimension of the fermions and the running of the coupling g. Thus the coupling will in reality become $m_\Psi \sim v[1 + \mathcal{O}(1)\frac{g^2}{4\pi^2}\ln(\frac{\phi}{v})]$, or after introducing the canonically normalized field φ from Eq. (4) and using $\xi \gg 1$,

$$m_\Psi \sim v\left(1 + \mathcal{O}(1)\frac{g^2}{4\pi^2}\frac{\varphi}{M_{\rm Pl}}\right). \quad (20)$$

This means that there will be Planck-suppressed couplings between φ and the fermions, and so the fermions will be produced directly by the Higgsflaton tachyonic preheating, and additional preheating stages as the field oscillates around the minimum [27,28].

Moreover, if there are fields with explicit mass terms in the theory, there will be mass-term induced direct couplings of φ to them, which are Planck suppressed, but may still be sufficiently large. This is most simply illustrated with an example of a scalar field χ defined by a Jordan-frame Lagrangian $\mathcal{L}_\chi = \sqrt{g}[-(\nabla\chi)^2/2 - U(\chi)]$. Upon changing to the Einstein frame metric variable, we find the leading order effective Lagrangian for χ,

$$\mathcal{L}_{\rm eff}(\hat{\chi}) = \sqrt{\hat{g}}\left[-\frac{1}{2}(\hat{\nabla}\chi)^2\left(\frac{v}{\phi_0}\right)^2 e^{-2(\varphi/M_{\rm Pl}\sqrt{6+1/\xi})} - \left(\frac{v}{\phi_0}\right)^4 e^{-4(\varphi/M_{\rm Pl}\sqrt{6+1/\xi})}U(\chi) + \ldots\right]. \quad (21)$$

If we expand this action in a series in φ around the minimum, where $\phi = \pm v$, and focus on the lowest order terms, we can see that the trilinear Lagrangian describing lowest order interactions is formed from keeping the kinetic term and the mass term for χ and the linear term in φ. This yields

$$\mathcal{L}_{\rm I} = \sqrt{\hat{g}}\frac{\varphi}{M_{\rm Pl}\sqrt{6+1/\xi}}[(\hat{\nabla}\chi)^2 + 2m_\chi^2\chi^2 + \ldots]. \quad (22)$$

Clearly, this trilinear term will yield the dominant channel for φ decay. The decay rate can now be calculated straightforwardly. Since one is interested at the decay of wave packets much smaller than the Hubble length, one can ignore the expansion of the universe and go to the locally Lorentzian frame, by replacing the metric in (22) by the Minkowski metric. Then since the leading order process is $\varphi \rightarrow 2\chi$, one can go to the momentum picture and evaluate (22) on shell. That yields $\mathcal{L}_{\rm I} = \frac{2m_\chi^2 - p_1\cdot p_2}{M_{\rm Pl}\sqrt{6+1/\xi}}\varphi\chi^2$, where p_k are the 4-momenta of the decay products, and in the center-of-mass (CM) frame of φ it reduces to, by recalling our metric signature to be $- + ++$ and using energy momentum conservation that yields $-p_1\cdot p_2 = \frac{m_\varphi^2}{2} - m_\chi^2$, an effective trilinear interaction

$$\mathcal{L}_{\rm I} = \frac{m_\chi^2 + m_\varphi^2/2}{M_{\rm Pl}\sqrt{6+1/\xi}}\varphi\chi^2, \quad (23)$$

which is just the standard scalar Yukawa term with the coupling constant $g = \frac{m_\chi^2 + m_\varphi^2/2}{M_{\rm Pl}\sqrt{6+1/\xi}}$. Therefore the decay rate $\varphi \rightarrow 2\chi$ is[8]

$$\Gamma_{\varphi\rightarrow 2\chi} = \frac{g^2}{8\pi m_\varphi}\sqrt{1 - 4\frac{m_\chi^2}{m_\varphi^2}} = \frac{(m_\chi^2 + m_\varphi^2/2)^2}{8\pi(6+1/\xi)M_{\rm Pl}^2 m_\varphi}\sqrt{1 - 4\frac{m_\chi^2}{m_\varphi^2}}. \quad (24)$$

Thus the gravitational decay time[9] when $m_{\hat{\chi}} \simeq m_\varphi$ and $\xi > 1$ is $\tau_{\rm gd} \sim M_{\rm Pl}^2 m_\varphi/m_\chi^4$ [20]. These extra channels will enhance reheating, and raise the reheating temperature: e.g. if $m_{\hat{\chi}} \simeq m_\varphi \simeq 3\times 10^{13}$ GeV, then $T_{\rm RH} \sim 10^8$ GeV. The reheating temperature of this range can directly be measured by observation of future space-based laser interferometers [29]. Moreover, in the presence of additional particles lighter than φ new channels will appear, enhancing $\Gamma \rightarrow \Gamma_{\rm total} = \sum_k \Gamma_k$. In any case, the Higgsflaton will settle down into the minimum rather efficiently. This is in fact good, because if any energy in it survived, it could overclose the universe. At any rate, this shows that the decay of the Higgsflaton would be efficient, and will convert the vacuumlike energy density of the Higgsflaton sector into normal particles. The precise details would of course depend on the exact structure of the physics which completes the standard model.

In lieu of a conclusion, let us state here that much of the dynamics presented here will remain a possibility for inflation even if LHC discovers the Higgs. In that case, however, many more theories involving light scalars may be plausible, and when identifying which may be the *raison d'etre* behind the inflaton, one may fall victim to a "tyranny of small decisions." The absence of the Higgs could, at least in this sense, help reduce the number of options for what lurks beyond the standard model, and point to a high scale inflation, that could be tested in the future searches for the primordial gravitational waves.

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[8] We are correcting here a typo involving a sign in the formula for $g \sim m_\chi^2 + m_\varphi^2/2$ in [20].

[9] The suppression of the decay rate by $M_{\rm Pl}$ arises since, by Eq. (4), the canonically normalized field φ is multiplied by $M_{\rm Pl}^{-1}$. This agrees with the revised version of [26].

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