On the Naturalness of Higgs Inflation

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Abstract

We critically examine the recent claim that the Standard Model Higgs boson $H$ could drive inflation in agreement with observations if $|H|^2$ has a strong coupling $\xi \sim 10^4$ to the Ricci curvature scalar. We first show that the effective theory approach upon which that claim is based ceases to be valid beyond a cutoff scale $\Lambda = m_p/\xi$, where $m_p$ is the reduced Planck mass. We then argue that knowing the Higgs potential profile for the field values relevant for inflation ($|H| > m_p/\sqrt{\xi} \gg \Lambda$) requires knowledge of the ultraviolet completion of the SM beyond $\Lambda$. In absence of such microscopic theory, the extrapolation of the pure SM potential beyond $\Lambda$ is unwarranted and the scenario is akin to other ad-hoc inflaton potentials afflicted with significant fine-tuning. The appealing naturalness of this minimal proposal is therefore lost.
1 The Standard Model Higgs as the Inflaton

Following [1] let us consider the SM Higgs sector corrected by a curvature-induced mass term for the Higgs boson:

\[
\mathcal{L}_{\text{Jordan}} = \frac{1}{2} m_p^2 R + \frac{1}{2} \xi h^2 R - \frac{1}{2} g^{\mu \nu} \partial_\mu h \partial_\nu h - \frac{1}{4!} \lambda (h^2 - v^2)^2 ,
\]

(1)

where \( m_p = 2.44 \times 10^{18} \text{ GeV} \) is the reduced Planck mass. Here \( h \) stands for the real neutral component of the Higgs doublet \( H \) that remains after the Higgs mechanism, the other scalar degrees of freedom in \( H \) being absorbed into the longitudinal components of the (omitted) gauge fields. This parametrization is natural at low energies, and much less so at energies well above the symmetry breaking scale, \( v = 246.22 \text{ GeV} \). Nevertheless, discussions of inflation often incorporate the requirement of a single-component inflaton field for phenomenological reasons. Therefore, we will assume that focusing on a single real component of the Higgs field is appropriate for our discussion at all energies of interest.

The effective Lagrangian (1) is assumed to be valid from the SM scale \( v \) up to some threshold below the reduced Planck mass \( m_p \). While the curvature coupling is irrelevant at low energies, it has quite remarkable consequences at large values of the Higgs field [2]. A suitable Weyl transformation of the metric

\[
g_{\mu \nu} \rightarrow (1 + \xi h^2 / m_p^2)^{-1} g_{\mu \nu} ,
\]

decouples the Higgs field from the curvature operator and reveals a new potential

\[
U(h) = \frac{\lambda}{4!} \frac{(h^2 - v^2)^2}{(1 + \xi h^2 / m_p^2)^2},
\]

(3)

with a modified large-field regime, capable of sustaining slow-roll inflation. In particular, one finds a plateau for \( h \gg m_p / \sqrt{\xi} \), with energy density \( m_p^4 / \xi^2 \). The phenomenological inflationary constraints (small slow-roll parameters and the right amplitude of density perturbations) can be met simply by choosing \( \xi \sim 10^4 \) [1] (see also [3, 4, 5] for further studies along these lines and [6] for some general effective-theory approaches to inflation).

These results suggest that the basic energy scale controlling inflation is \( \Lambda_I \equiv m_p / \sqrt{\xi} \). Since \( \xi \) is large, we are guaranteed that Planck-suppressed effective operators have negligible effect on the effective potential. In any case, the emergence of the plateau at large field strength is a remarkable feature that is not at all expected at the level of the innocent-looking Lagrangian in the so-called Jordan frame [1]. While the occurrence of inflationary plateaus is often attributed to peculiar dynamical features or plain fine-tuning, we seem to be producing one here, just from a quite generic low-energy effective Lagrangian.

In this note we make two simple observations. First, we show in section 2 that the actual effective cutoff of [1] is \( \Lambda = m_p / \xi \), i.e. much lower than the energy scales determining the properties of the plateau itself. This suggests that, despite appearances, there is a strong fine-tuning implicit in [3]. In connection with this, our second observation (section
3) is that the plateau is extremely sensitive to the presence of operators suppressed by the scale $\Lambda$ which are generically expected to appear. The claimed naturalness of the scenario is therefore lost.

## 2 Effective Cutoff Scale

The presence of nonrenormalizable operators in the effective Lagrangian gives an upper bound on its own cutoff scale. In the weak-field regime, the operator $h^2 R$ has dimension five by power counting. Expanding around Minkowski space with a canonically normalized graviton $g_{\mu\nu} = \eta_{\mu\nu} + m_p^{-1} \gamma_{\mu\nu}$ we find a leading term

$$\frac{\xi}{m_p} h^2 \eta^{\mu\nu} \partial^2 \gamma_{\mu\nu} + \ldots$$

where the dots stand for the other tensor structures conforming the linearized approximation of the Ricci scalar. The contribution of $\frac{\xi}{m_p} h^2$ to tree processes is controlled by the ratio $\frac{\xi E}{m_p}$ at typical energy $E$, becoming strongly coupled at $E \sim \Lambda = m_p/\xi$. Hence the effective cutoff of the nonrenormalizable Lagrangian is $\Lambda = m_p/\xi$. (Other higher-order operators containing two or more gravitons are suppressed instead by $\Lambda_I = m_p/\sqrt{\xi}$ which is a much higher energy scale for $\xi \gg 1$.)

The same effective cutoff can be obtained in the Einstein frame, in which all the non-linearities are transferred to the scalar sector. Both the Einstein-frame potential and the rescaled kinetic term of the Higgs field only contain higher-dimensional operators suppressed by the naive inflation scale $m_p/\sqrt{\xi}$. There is, however, an extra contribution to the two-derivative effective action coming from the Weyl rescaling of the curvature, i.e.

$$-\frac{3\xi^2}{m_p^2} \frac{h^2}{(1 + \xi h^2/m_p^2)^2} (\partial h)^2$$

whose leading term is the dimension-six operator

$$-3\frac{\xi^2}{m_p^2} h^2 (\partial h)^2,$$

with effective cutoff $\Lambda = m_p/\xi$.

All these nonlinearities of the Higgs kinetic term can be hidden by the field redefinition

$$d\phi = dh \left[1 + 6\xi^2 (h/m_p)^2 + \xi (h/m_p)^2 \right]^{1/2}$$

leading to a canonically normalized scalar sector in the Einstein frame,

$$\mathcal{L}_{\text{Einstein}} = \frac{1}{2} m_p^2 R - \frac{1}{2} (\partial \phi)^2 - U(\phi),$$

1In a non-zero Higgs background the normalization of the graviton involves the modified Planck mass $m_p^2 + \xi \langle h \rangle^2$. This does not affect qualitatively our results. We also assume $\xi > 0$, since $\xi < 0$ lowers the effective Planck mass at large $\langle h \rangle$, making gravity effects even stronger.
where the new potential \( U(\phi) \) is given by (3), after we solve the field redefinition (7) in favor of the canonical Higgs field \( \phi \) in the Einstein frame. All couplings appearing in \( U(\phi) \) are non-redundant, since we have exhausted the freedom in nonlinear field redefinitions. Working at small values of the Higgs field we find

\[
h = \phi \left[ 1 - (\xi \phi / m_p)^2 \right] + \text{higher order terms.}
\]

Substituting (9) into a generic \( h^4 \) term of the potential, the factor of \( 1 - \xi^2 \phi^2 / m_p^2 \) induces a dimension-six effective operator proportional to

\[
\frac{\xi^2}{m_p^2} \phi^6,
\]

again showing an effective cutoff \( \Lambda = m_p / \xi \) (we neglect \( O(1) \) factors like \( \sqrt{\Lambda} \)). The same cutoff appears after substituting (9) in Higgs Yukawa couplings like that of the top quark.\(^2\) It is thus reassuring that the same high energy cutoff appears in both frames, as corresponds to a meaningful physical scale [7]. The discussion of effective operators in the Einstein frame has the advantage of decoupling the problem from the possible subtleties related to the diffeomorphism invariance in the gravity sector.

We have confirmed that the natural cutoff of the effective Lagrangian expanded around the SM region \( \phi \sim v \) is given by \( \Lambda \) above. In fact, this piece of limiting information is nearly all we know about the effective action extrapolated to higher energy scales. Can we trust the effective potential for the field values \( h > \Lambda_I = m_p / \sqrt{\xi} \) relevant for inflation? One could argue that quantum fluctuations around that large Higgs background are very weakly coupled and we do not see there any strong coupling effect. However, the very existence of the plateau, at energy density \( m_p^4 / \xi^2 = \Lambda_I^4 \gg \Lambda^4 \) is suspect, as we discuss next.

3 UV Sensitivity of the Inflationary Plateau

The field redefinitions carried out so far convert the simple Lagrangian (11) into a non-polynomial scalar model in terms of the canonical field \( \phi \). From the point of view of effective field theory this non-polynomial potential, \( U(\phi) \), is quite fine-tuned though, since all its terms depend upon a single dimensionless parameter \( \xi \), apart from the overall

\(^2\) To clarify further the point already made in our published paper let us stress that the operator (10) is the lowest one (in the scalar sector) revealing the true cutoff, after all ambiguities due to redundant couplings are removed. In this respect, a scattering process like \( \phi \phi \to \phi \phi \phi \phi \) is more appropriate to discuss the breakdown of the effective theory than \( \phi \phi \to \phi \phi \). Other dimension-6 operators playing a similar role are \( h_I^2 \phi^3 \bar{\psi} \psi / \Lambda^2 \) or \( g^2 \phi^4 W_\mu^+ W^\mu^- / \Lambda^2 \). As a general comment, for the cutoff to be \( \Lambda = M_p / \xi \), it is crucial that the Higgs scalar has couplings to other sectors or to itself (or the Goldstone components of the Higgs doublet). If all those couplings were zero (and the Goldstones were absent) the theory in the Einstein frame would correspond to a free scalar minimally coupled to gravity (and in the Jordan frame that would cause miraculous cancellations in the scattering amplitudes of \( h \)). We thank G. Giudice and A. Riotto for a discussion on this point.
constant $\lambda$. It is then interesting to study the sensitivity of the plateau at $\phi > m_p/\sqrt{\xi}$ to the presence of higher-dimensional operators with the same basic structure of the Jordan frame Lagrangian (I) in the gravitational sector.

Higher order corrections to the scalar potential in the Jordan frame have the form

$$V(h) = \sum_{n \geq 4} \frac{1}{\Lambda^{2n-4}} \lambda_n(h) h^{2n}, \quad (11)$$

where the effective couplings $\lambda_n$ may depend on $\log(h/\Lambda)$, a mild field-dependence induced by radiative corrections to the bare potential. In general, in the presence of the coupling (4) all terms in the expansion (11) will be generated by radiative corrections, even if they were not present in the Wilsonian effective potential at the cutoff scale $\Lambda$. Hence, the low-energy expansion in the vicinity of SM field strengths suggests that the generic potential has the form

$$V(h) = \Lambda^4 V(h/\Lambda), \quad (12)$$

with $V(x)$ an even function of $O(1)$, with a typical scale of variation of $O(1)$.

Analogous considerations for the curvature term lead to a tower of operators in the linearized gravity approximation,

$$\sum_{n \geq 2} \frac{1}{\Lambda^{2n-1}} \alpha_n(h) h^{2n} \eta^{\mu\nu} \partial^2 \gamma_{\mu\nu} + \ldots, \quad (13)$$

where $\alpha_n(h)$ will also have logarithmic field dependence, matched at the threshold scale $\Lambda$ to the high energy theory. Using $R \sim m_p^{-1} \partial^2 \gamma$, we conclude that the generic curvature coupling at leading order in the Ricci scalar has the form

$$\frac{1}{2} m_p \Lambda f(h/\Lambda) R, \quad (14)$$

where, again $f(x)$ is an even function of $O(1)$, with typical scale of variation of $O(1)$. While the full effective action below the scale $\Lambda$ is certainly not exhausted by the functions $V(x)$ and $f(x)$, our qualitative points can be put forward by concentrating on those terms directly involved in the inflationary mechanism under discussion.

The Weyl transformation decoupling the curvature from the scalar field is now

$$g_{\mu\nu} \to \left[ 1 + \frac{\Lambda}{m_p} f(h/\Lambda) \right]^{-1} g_{\mu\nu}, \quad (15)$$

and the scalar field redefinition that achieves canonical normalization in the Einstein frame:

$$d\phi = dh \left[ 1 + \xi^{-1} f(h/\Lambda) + (3/2)[f'(h/\Lambda)]^2 \right]^{1/2} \frac{1}{1 + \xi^{-1} f(h/\Lambda)}, \quad (16)$$

3We assume here a mass-independent renormalization scheme in the effective field theory calculations.
4The generic mass for $h$ is $O(\Lambda)$, a manifestation of the hierarchy problem, that we do not address here.
where \( f'(x) \equiv df(x)/dx \) and we have defined \( \xi \equiv m_p/\Lambda \). The corresponding scalar potential takes the form

\[
U(\phi) = \frac{\Lambda^4 V(h(\phi)/\Lambda)}{\left[1 + \frac{\Lambda}{m_p} f(h(\phi)/\Lambda)\right]^2} ;
\]

(17)

where \( h(\phi) \) is the solution of the field redefinition (16).

We can now discuss the general conditions for the emergence of an inflationary plateau in the Einstein-frame potential \( U(\phi) \). Going back to (16), there are two regimes depending on whether \( f \ll \xi \) or \( f \gg \xi \). In the first case the dependence on \( \xi \) drops out and we have a generalization of (9), i.e. \( h \sim \phi[1 + \mathcal{O}(\phi^2/\Lambda^2)] \), the weak-field regime that exposes \( \Lambda \) as the relevant scale of the problem. In this regime, any plateau must be explicitly engineered in the potential function \( V(h) \) from the beginning.

On the other hand, any unexpected plateau extending over the region \( f \gg \xi \) will have the form

\[
U(\phi) \approx \Lambda_I^4 \frac{V(h(\phi)/\Lambda)}{[f(h(\phi)/\Lambda)]^2} ;
\]

(18)

where \( \Lambda_I = m_p/\sqrt{\xi} \) is the energy scale of the plateau. Hence, a flat potential at large values of the Higgs field requires a functional constraint \( V(h/\Lambda) \approx [f(h/\Lambda)]^2 \) over a sufficiently large region so that we sustain enough slow roll. Since the functions \( V(x) \) and \( f(x) \) are \textit{a priori} independent, this constraint is equivalent to the explicit tuning of one function.

We can be slightly more specific if we make the extra technical assumption that \( f'(x)^2 \) is not anomalously small compared to \( f(x) \) at large \( x \) (which includes for instance polynomials or trigonometric polynomials). Then we may approximate the field redefinition (16) as

\[
\frac{d\phi}{m_p} \sim \sqrt{\frac{3}{2}} \frac{f'(h/\Lambda)}{f(h/\Lambda)} d(h/\Lambda) = \sqrt{\frac{3}{2}} d(\log f) .
\]

(19)

Hence, we find

\[
f \sim \xi \exp\left(\frac{a \phi}{m_p}\right) \quad \text{for} \quad f \gg \xi ,
\]

(20)

where we have defined \( a = \sqrt{2/3} \) and we have chosen the additive normalization of \( \phi \) so that the matching between the two regimes, at \( f \sim \xi \), corresponds to \( \phi \sim m_p \).

Let us consider the function \( V(x)/[f(x)]^2 \) and eliminate the \( x \) variable in favor of the \( f \) variable, i.e. we invert \( x = x(f) \) and write

\[
\frac{V(x)}{[f(x)]^2} = \frac{V(x(f))}{f^2} \equiv F(f) .
\]

(21)

We know that near the origin of field space, and neglecting low-energy scales of the SM, \( V \ll \Lambda^4 \), \( F(f) \) is a polynomial function with \( \mathcal{O}(1) \) coefficients. In the case of the Lagrangian
we have $F(f) = 1$ exactly. On the other hand, possible plateaus have to do with the large $f$ behavior of the function $F(f)$. Any plateau at large $f$ that ends around $f \sim \xi$ is described by a function that admits a large $f$ expansion of the form

$$F(f) \sim 1 + \mathcal{O}(\xi/f) ,$$

with $O(1)$ coefficients in general. On the other hand, the complete potential (17) reads

$$U = \Lambda^4 I_4 F(f) (1 + \frac{\xi}{f})^2 = \Lambda^4 I_4 [1 + \mathcal{O}(\xi/f)] = \Lambda^4 I_4 \left[ 1 + \mathcal{O}\left( e^{-a\phi/m_p} \right) \right] ,$$

in the plateau regime. Hence, we have seen that the general picture for Higgs inflation will come to fruition whenever we have a function $F(f)$ which approaches unity with power-like accuracy in the region $f \gg \xi$. The phenomenological conditions for inflation will be satisfied just as in ref. [1], with $O(1)$ coefficients in all the power expansions. One simply needs to adjust the single large parameter $\xi \sim 10^4$. So, where is the fine tuning?

The fine-tuning is of course in the assumption that the function $F(f)$ approaches unity at large $f$. Notice that, in general, $F(f)$ cannot even be defined, because unless $f(x)$ is monotonic, we will not be able to solve $x$ as a function of $f$ in the first place. For instance, if $\mathcal{V}(x)$ and $f(x)$ are ‘landscape-like’ with a succession of maxima and minima separated by a distance $\mathcal{O}(1)$ in $x$, we will not have a plateau in terms of the canonical field. This is very clear in the particular example in which we take $f(x) = x^2$ as in ref. [1], and $\mathcal{V}(x) = \mathcal{P}(x^2)$ is a pseudoperiodic function with approximate period of $\mathcal{O}(1)$. The asymptotic form of the potential is then

$$U \rightarrow \Lambda^4 e^{-2a\phi/m_p} \mathcal{P} \left( \xi e^{a\phi/m_p} \right) ,$$

We see that, in $\phi$ space, we have a singular function with ever increasing oscillation rate, which in addition is quenched by a decaying exponential, hardly a good environment for inflation!

From a physical point of view, all we can expect to know is the small field expansion of the various functions involved, $\mathcal{V}(x)$ and $f(x)$, collected from fits to amplitudes measured at energies $E \ll \Lambda$. From these functions we obtain $F(f)$, which in this regime looks like a polynomial in $f$. Suppose we have determined that it has the form

$$F(f) = 1 + \sum_{n \geq 1} \alpha_n f^n ,$$

from fits to low-energy data (fixing the coefficient of $n$-th order requires measuring effective operators of dimension $2n$ in the Higgs sector). The so-called Higgs inflation will take place if this function approaches unity with power-like accuracy as $f > \xi$. Of course, since $\xi \gg 1$, there is no way we can infer the behavior at $f/\xi > 1$ from the above low-energy power expansion. If we make an arbitrary truncation at $n \leq N$ (this will always be the case
because the experimental accuracy will be finite) and extrapolate the result to large $f$, we either need miraculous cancellations, or else we must bound

$$|\alpha_n| \ll \frac{1}{\xi^n}, \quad (26)$$

for all $n \leq N$, which is a pretty strong fine-tuning indeed. In other words, if one starts with a generic potential without plateau in the Einstein frame, the potential transformed back into the Jordan frame through eqs. (2) and (7) will look at low-energy like a normal SM Higgs potential plus a tower of non-renormalizable operators difficult to measure. Without detailed information about the structure of that tower of operators one cannot predict the presence of a plateau in the Einstein frame.

It is interesting to reinterpret this fine-tuning in terms of scales again. Let us set $f(x) = x^2$, so that (25) translates into an expansion of the dimensionless potential function:

$$V = f^2 + \sum_{n \geq 1} \alpha_n f^{n+2} = x^4 + \sum_{n \geq 1} \alpha_n x^{2n+4}, \quad (27)$$

which in turn translates into the following expansion of the Jordan-frame potential

$$V(h) = h^4 + \sum_{n \geq 1} \alpha_n \frac{h^{2n+4}}{\Lambda^{2n}} = h^4 + \sum_{n \geq 1} \beta_n \frac{h^{2n+4}}{\Lambda^2}, \quad (28)$$

where we have written $\beta_n = \alpha_n \xi^n$. The very stringent fine-tuning condition (26) actually means that the potential is ‘natural’ with respect to the higher scale $\Lambda_I = \sqrt{\xi} \Lambda$, with at most a slow-roll fine-tuning\(^5\) of the coefficients $|\beta_n| \ll 1$.

Therefore, in this example Higgs inflation is supported by the fact that, while the non-minimal curvature coupling in the Jordan frame is heralding the presence of an effective cutoff at the scale $\Lambda$, the effective scalar potential in the same frame has a hierarchically larger effective cutoff $\Lambda_I$. However, even this situation is not enough to guarantee the existence of the plateau, since any significant deviation from the behavior $f(x) = x^2$ at large values of $f \sim \xi$ will destroy the plateau as well.

We conclude that a very peculiar structure of the effective action is required, with a hierarchical separation of scales within different sectors, combined with a strong correlation of these sectors at field strengths in excess of the higher scale. Even if one such effective action is set at the cutoff scale $\Lambda$, radiative corrections from quantum loops below $\Lambda$ will generically upset such fragile structure. Needless to say, this problem afflicts also other inflationary models that use the same large coupling of the inflaton to the Ricci scalar as a way of obtaining a plateau without invoking any symmetry reason.

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\(^5\)By slow-roll fine-tuning we mean the ‘technical’ tuning needed to keep the slow-roll parameters small.
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Note added

While this paper was being drafted, ref. [8] appeared. In this work related issues are discussed, including the identification of a low effective cutoff scale $m_p/\xi$, and the possible pernicious influence of higher non-renormalizable Higgs operators (coming from integrating out some heavy sector coupled to the Higgs) on the potential plateau. While the authors of ref. [8] examine the applicability of the semiclassical approximation and leave open the door for a (narrow) parameter region where inflation could be viable, we emphasize here that the very existence of the inflationary plateau is compromised by the result $\Lambda = m_p/\xi$.

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