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Higgs–inflaton coupling from reheating and the metastable Universe

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ABSTRACT

Current Higgs boson and top quark data favor metastability of our vacuum which raises questions as to why the Universe has chosen an energetically disfavored state and remained there during inflation. In this Letter, we point out that these problems can be solved by a Higgs–inflaton coupling which appears in realistic models of inflation. Since an inflaton must couple to the Standard Model particles either directly or indirectly, such a coupling is generated radiatively, even if absent at tree level. As a result, the dynamics of the Higgs field can change dramatically.

The current Higgs mass $m_h = 125.15 \pm 0.24$ GeV and the top quark mass $m_t = 173.34 \pm 0.76 \pm 0.3$ GeV indicate that in the Standard Model (SM) the Higgs quartic coupling turns negative at high energies implying metastability of the electroweak (EW) vacuum at 99% CL [1]. The (much deeper) true minimum of the scalar potential appears to be at very large field values. In the cosmological context, this poses a pressing question why the Universe has chosen an energetically disfavored state and why it remained there during inflation despite quantum fluctuations.

In this Letter, we argue that these puzzles can be resolved by a Higgs–inflaton coupling [2] which appears in realistic models of inflation. Indeed, the energy transfer from the inflaton to the SM fields necessitates interaction between the two in some form. This in turn induces a Higgs–inflaton coupling via quantum effects, even if it is absent at tree level. We find that the loop induced coupling can be sufficiently large to make a crucial impact on the Higgs field evolution.

Another factor that can affect the Higgs field dynamics is the non-minimal scalar coupling to gravity, which creates an effective mass term for the Higgs field [3,4]. Here we assume such a coupling to be negligible. The effect of quantum fluctuations during inflation has recently been considered in [5,6]. The conclusion is that the Hubble rate H above the Higgs instability scale leads to destabilization of the EW vacuum, which poses a problem for this class of inflationary models. Related issues have been studied in [7–9].

The Higgs potential at large field values is approximated by [10]

$$V_h \simeq \frac{\lambda_h(h)}{4}\, h^4 \,, \tag{1}$$

where we have assumed the unitary gauge $H^T = (0, h/\sqrt{2})$ and $\lambda_h(h)$ is a logarithmic function of the Higgs field. The current data indicate that λ_h turns negative at around 10^{10} GeV [1], although the uncertainties are still significant. In the early Universe, the Higgs potential is modified by the Higgs–inflaton coupling $V_{h\phi}$ with the full scalar potential being

$$V = V_h + V_{h\phi} + V_\phi \,, \tag{2}$$

where V_ϕ is the inflaton potential. Since the inflaton must couple to the SM fields either directly or through mediators as required by successful reheating, quantum corrections induce a Higgs–inflaton interaction.

In what follows, we consider a few representative examples of reheating models. We focus on the Higgs couplings to the inflaton ϕ which are *required* by renormalizability of the model. Such couplings are induced radiatively with divergent coefficients and necessitate the corresponding counterterms. The dim-4 Higgs–inflaton interaction takes the form

$$V_{h\phi} = \frac{\lambda_{h\phi}}{4}\, h^2\phi^2 + \frac{\sigma_{h\phi}}{2}\, h^2\phi \,, \tag{3}$$

where $\lambda_{h\phi}$ and $\sigma_{h\phi}$ are model-dependent couplings. As we show below, the range of $\lambda_{h\phi}$ relevant to the Higgs potential stabilization is between 10^{-10} and 10^{-6} (see also [2]). For definiteness, we choose a quadratic inflaton potential [11] as a representative example of large field inflationary models,

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$$V_\phi = \frac{m^2}{2}\,\phi^2 + \Delta V_{1-\mathrm{loop}}\,, \tag{4}$$

where $m \simeq 10^{-5} M_{\mathrm{Pl}}$ and $\Delta V_{1-\mathrm{loop}}$ is the radiative correction generated by various couplings of the model. We require this correction to be sufficiently small such that the predictions for cosmological observables of the ϕ^2-model are not affected, although some quantum effects can be beneficial [12]. The divergent contributions to $\Delta V_{1-\mathrm{loop}}$ are renormalized in the usual fashion and the result is given by the Coleman–Weinberg potential [13]. The leading term at large ϕ is the quartic coupling

$$\Delta V_{1-\mathrm{loop}} \simeq \frac{\lambda_\phi(\phi)}{4}\,\phi^4\,, \tag{5}$$

with λ_ϕ being logarithmically dependent on ϕ.

The energy transfer from the inflaton to the SM fields in general proceeds both through non-perturbative effects and perturbative inflaton decay [14,15]. In what follows, we make the simplifying assumption that the reheating is dominated by the perturbative inflaton decay such that the reheating temperature is given by $T_R \simeq 0.2\sqrt{\Gamma M_{Pl}}$, where Γ is the inflaton decay rate. While this assumption is essential for establishing a correlation between $\lambda_{h\phi}$ and T_R, it does not affect the range of $\lambda_{h\phi}$ consistent with the inflationary predictions. We consider three representative reheating scenarios which assume no tree level interaction between the Higgs and the inflaton, and compute the consequent loop-induced couplings.

1. Reheating via right-handed neutrinos

The inflaton energy is transferred to the SM sector via its decay into right-handed Majorana neutrinos ν_R which in turn produce SM matter. The added benefit of this model is that the heavy neutrinos may also be responsible for the matter–antimatter asymmetry of the Universe via leptogenesis [16]. The relevant tree level Lagrangian reads

$$-\Delta\mathcal{L} = \frac{\lambda_\nu}{2}\,\phi\nu_R\nu_R + y_\nu\,\bar{l}_L\cdot H^*\,\nu_R + \frac{M}{2}\,\nu_R\nu_R + \mathrm{h.c.}\,, \tag{6}$$

where l_L is the lepton doublet, M is chosen to be real and we have assumed that a single ν_R species dominates. These interactions generate a coupling between the Higgs and the inflation at 1 loop (Fig. 1). Since we are interested in the size of the radiatively induced couplings, let us impose the renormalization condition that they vanish at a given high energy scale, say the Planck scale $M_{\mathrm{Pl}} = 2.4 \times 10^{18}$ GeV. Then, a finite correction is induced at the scale relevant to the inflationary dynamics, which we take to be the Hubble rate $H = m\phi/\left(\sqrt{6}M_{\mathrm{Pl}}\right)$, with other choices leading to similar results. We find in the leading-log approximation,

$$\begin{aligned}
\lambda_{h\phi} &\simeq \frac{|\lambda_\nu y_\nu|^2}{2\pi^2}\ln\frac{M_{\mathrm{Pl}}}{H}\,,\\
\sigma_{h\phi} &\simeq -\frac{M|y_\nu|^2\mathrm{Re}\lambda_\nu}{2\pi^2}\ln\frac{M_{\mathrm{Pl}}}{H}\,,\\
\lambda_\phi &\simeq \frac{|\lambda_\nu|^4}{4\pi^2}\ln\frac{M_{\mathrm{Pl}}}{H}\,.
\end{aligned} \tag{7}$$

Here we have chosen the same renormalization condition for λ_ϕ and $\lambda_{h\phi}$, $\sigma_{h\phi}$. Since the dependence on the renormalization scale is only logarithmic, this assumption does not affect our results. The most important constraint on the couplings is imposed by the inflationary predictions. Requiring $\lambda_\phi\phi^4/4 \ll m^2\phi^2/2$ in the last 60 e-folds of expansion (see e.g. [17]), we find $\lambda_\phi \ll 2\times 10^{-12}$

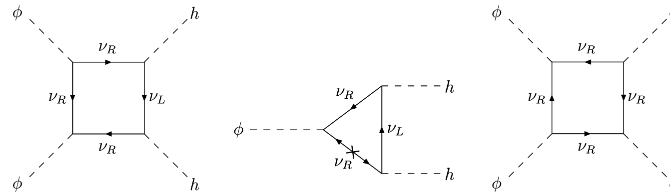

Fig. 1. Leading radiatively induced scalar couplings via the right-handed neutrinos. (Diagrams with the same topology are not shown.)

and therefore $\lambda_\nu < 1\times 10^{-3}$. The seesaw mechanism also limits the size of the Yukawa coupling y_ν. The experimental constraints on the mass of the active neutrinos require approximately $(y_\nu v)^2/M < 1$ eV. Assuming that the perturbative decay of the inflaton dominates, the mass of the right-handed neutrinos is bounded by $M < 10^{13}$ GeV, which in turn implies $y_\nu < 0.6$. We therefore get an upper bound on the size of the Higgs–inflaton coupling,

$$\lambda_{h\phi} < 2\times 10^{-7}\,. \tag{8}$$

Note that $\lambda_{h\phi}$ is positive and thus the inflaton creates a positive effective mass term for the Higgs. The trilinear ϕh^2 term is irrelevant as long as $|\lambda_\nu|\phi \gg M$, which is the case for all interesting applications. (Similarly, the cubic term ϕ^3 is negligible.)

During the inflaton oscillation stage, the magnitude of ϕ decreases as $1/t$. When the effective masses of ν_R and h turn sufficiently small, the decays $\phi\to\nu_R\nu_R$, $\phi\to hh$ become allowed. The constraints above imply $\Gamma(\phi\to\nu_R\nu_R) \gg \Gamma(\phi\to hh)$ and therefore the total inflaton decay width is $\Gamma = \frac{|\lambda_\nu|^2}{32\pi}\,m$, where we have neglected the ν_R mass compared to that of the inflaton. Assuming that the right-handed neutrinos decay promptly and the products thermalize (or ν_R themselves thermalize) so that $T_R \simeq 0.2\sqrt{\Gamma M_{\mathrm{Pl}}}$, we find the following correlation between the Higgs–inflaton coupling and the reheating temperature T_R,

$$\lambda_{h\phi} \simeq 5\times 10^{-7}\,|y_\nu|^2\left(\frac{T_R}{1.5\times 10^{11}\ \mathrm{GeV}}\right)^2\,, \tag{9}$$

where T_R is bounded by 1.5×10^{11} GeV. Note that this relation holds only under the assumption of perturbative reheating. Therefore, for the neutrino Yukawa coupling and the reheating temperature within one–two orders of magnitude from their upper bounds, the dynamics of the Higgs evolution change drastically. Similar conclusions apply to models with multiple ν_R species.

2. Reheating and non-renormalizable operators

A common approach to reheating is to assume the presence of non-renormalizable operators that couple the inflaton to the SM fields. Let us consider a representative example of the following operators

$$O_1 = \frac{1}{\Lambda_1}\phi\,\bar{q}_L\cdot H^*t_R\,,\quad O_2 = \frac{1}{\Lambda_2}\phi\,G_{\mu\nu}G^{\mu\nu}\,, \tag{10}$$

where $\Lambda_{1,2}$ are some scales, $G_{\mu\nu}$ is the gluon field strength and q_L, t_R are the third generation quarks. These couplings allow for a direct decay of the inflaton into the SM particles. It is again clear that a Higgs–inflaton interaction is induced radiatively. In order to calculate the 1-loop couplings reliably, one needs to complete the model in the ultraviolet (UV). The simplest possibility to obtain an effective dim-5 operator is to integrate out a heavy fermion. Therefore, we introduce vector-like quarks Q_L, Q_R with the tree level interactions

$$-\Delta\mathcal{L} = y_Q\,\bar{q}_L\cdot H^*\,Q_R + \lambda_Q\,\phi\,\bar{Q}_L t_R + \mathcal{M}\,\bar{Q}_L Q_R + \mathrm{h.c.}\,, \tag{11}$$