Cosmological Perturbations from the Standard Model Higgs

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Abstract

We propose that the Standard Model (SM) Higgs is responsible for generating the cosmological perturbations of the universe by acting as an isocurvature mode during a de Sitter inflationary stage. In view of the recent ATLAS and CMS results for the Higgs mass, this can happen if the Hubble rate during inflation is in the range $(10^{10} - 10^{14})$ GeV (depending on the SM parameters). Implications for the detection of primordial tensor perturbations through the $B$-mode of CMB polarization via the PLANCK satellite are discussed. For example, if the Higgs mass value is confirmed to be $m_h = 125.5$ GeV and $m_t, \alpha_s$ are at their central values, our mechanism predicts tensor perturbations too small to be detected in the near future. On the other hand, if tensor perturbations will be detected by PLANCK through the $B$-mode of CMB, then there is a definite relation between the Higgs and top masses, making the mechanism predictive and falsifiable.
1 Introduction

Experimental data recently reported by ATLAS [1, 2] and CMS [3, 4] are consistent with the discovery of the Standard Model (SM) Higgs boson, with mass around 125 – 126 GeV.

Such a light Higgs is in good agreement with the indirect indications derived from electroweak precisions constraints [5] under the hypothesis of negligible contributions of physics beyond the SM. Moreover, no clear signal of non-SM physics has emerged yet from collider searches.

Motivated by this experimental situation, in the present paper we try to answer a simple question: can the SM Higgs be responsible for the cosmological perturbations we observe in the universe? As we will see, the answer is: yes.

One of the basic ideas of modern cosmology is that there was an epoch early in the history of the universe when potential, or vacuum, energy associated to a scalar field, the inflaton, dominated other forms of energy density such as matter or radiation. During such a vacuum-dominated era the scale factor grew (nearly) exponentially in time. During this phase, dubbed inflation [6–8], a small, smooth spatial region of size less than the Hubble radius could grow so large as to easily encompass the comoving volume of the entire presently observable universe. If the universe underwent such a period of rapid expansion, one can understand why the observed universe is so homogeneous and isotropic to high accuracy.

Inflation has also become the dominant paradigm for understanding the initial conditions for the Large Scale Structure (LSS) formation and for Cosmic Microwave Background (CMB) anisotropy. In the inflationary picture, primordial density and gravity-wave fluctuations are created from quantum fluctuations “redshifted” out of the horizon during an early period of superluminal expansion of the universe, where they are “frozen” [9–14]. Perturbations at the surface of last scattering are observable as temperature anisotropy in the CMB. The last and most impressive confirmation of the inflationary paradigm has been recently provided by the data of the Wilkinson Microwave Anistropy Probe (WMAP) mission which has marked the beginning of the precision era of the CMB measurements in space [15].

Despite the simplicity of the inflationary paradigm, the mechanism by which cosmological adiabatic perturbations are generated is not yet fully established. In the standard picture, the observed density perturbations are due to fluctuations of the inflaton field itself. When inflation ends, the inflaton oscillates about the minimum of its potential and decays, thereby reheating the universe. As a result of the fluctuations each region of the universe goes through the same history but at slightly different times. The final temperature anisotropies are caused by the fact that inflation lasts different amounts of time in different regions of the universe leading to adiabatic perturbations.

Can the SM Higgs and its potential be responsible for inflation and, at the same time, the
generation of anisotropies?

The answer is: most probably, not both [16]. The basic problem is that the requirement of having enough e-folds of inflation requires the SM potential to be flat enough, but this conflicts with the requirement that quantum fluctuation of the Higgs inflaton should also generate the observed power spectrum of anisotropies. Indeed the height of the SM potential in its flat region is predicted and cannot be arbitrarily adjusted to be as low as needed. This problem can be, in principle, solved by non-minimal coupling of the SM Higgs scalar field $h$ to the Ricci scalar $R$ of the form $\xi h^2 R$ [17–19]. The effect of this interaction is to flatten the Higgs potential (or any other potential) above the scale $M_{Pl}/\sqrt{\xi}$ providing a platform for slow-roll inflation. A correct normalization of the spectrum of primordial fluctuations fixes the value of the coupling constant $\xi$ to be larger than about $10^4$.

This minimal inflationary scenario faces though a couple of issues. First, perturbative unitarity is violated at some scale lower than $M_{Pl}$ [20] (see however Ref. [21]), possibly implying the presence of new degrees of freedom which may change the inflationary dynamics. Secondly, the scenario requires stability of the potential up to the inflationary scale $M_{Pl}/\sqrt{\xi}$. With the LHC indication of a Higgs mass in the range $m_h = (125–126)$ GeV, this simplest version of Higgs inflation is disfavored as the Higgs potential develops an instability at much smaller scales, unless the top mass is below $\sim 171$ GeV [22–25].

These considerations lead us to believe that the SM Higgs field may not be responsible for both driving inflation and generating the cosmological perturbations at the same time. Let us therefore be more modest and drop off the requirement that the SM Higgs potential was responsible for inflation. Our goal is to show that the SM Higgs can nevertheless play a role in giving rise to the LSS and CMB anisotropies. Indeed, the standard scenario for the generation of the perturbations, where it is the same scalar field to drive inflation and to source the perturbations, is not the only option. For instance, an alternative to the standard scenario is represented by the curvaton mechanism [26–29] where the final curvature perturbation $\zeta$ is produced from an initial isocurvature perturbation associated to the quantum fluctuations of a light scalar field (other than the inflaton), the curvaton, whose energy density is negligible during inflation. The curvaton isocurvature perturbations are transformed into adiabatic ones when the curvaton decays into radiation much after the end of inflation. Alternatives to the curvaton model are those models characterized by the curvature perturbation being generated by an inhomogeneity in the decay rate [30] or the mass [31] of the particles responsible for the reheating after inflation. Other opportunities for generating the curvature perturbation occur at the end of inflation [32] and during preheating [33].

All these alternative models to generate the cosmological perturbations have their strength in the fact that all scalar fields during a period of de Sitter with a mass smaller than the Hubble rate $H$ during inflation are inevitably quantum-mechanically excited with a final superhorizon flat spectrum. Furthermore, they have in common that the comoving curvature perturbation is
generated on superhorizon scales when the isocurvature perturbation, which is associated to the fluctuations of these light scalar fields different from the inflaton, is converted into curvature perturbation after (or at the end) of inflation.

In the rest of the paper we will therefore assume that there was an inflationary period of accelerated expansion during the primordial evolutionary stage of the universe. This de Sitter period, induced by some unspecified vacuum energy, is characterized by a Hubble rate $H$. We will also assume that the perturbations are generated through the Higgs field. This will allow us to play with two independent parameters, the SM Higgs mass $m_h$ and the Hubble rate $H$. Also, for simplicity, we assume that the inflaton sector does not alter the SM sector, e.g. the Higgs potential (see also Ref. [34] for an alternative idea where the Higgs sector of the SM is minimally coupled to asymptotically safe gravity).

Our considerations will have direct observational consequences once one realizes that the Hubble rate parametrizes the amount of tensor perturbations during inflation. During the inflationary epoch, tensor perturbations, as for any other massless scalar field, are quantum-mechanically generated. They can give rise to $B$-modes of polarization of the CMB radiation through Thomson scatterings of the CMB photons off free electrons at last scattering [35]. The amplitude of the $B$-modes depends on the amplitude of the gravity waves generated during inflation, which in turn depends on the energy scale at which inflation occurred. The tensor-to-scalar power ratio is given by $T/S \approx (H/3.0 \times 10^{14}\text{ GeV})^2$. Current CMB anisotropy data impose the upper bound $T/S \lesssim 0.24$ [36]. The possibility of detecting gravity waves from inflation via $B$-modes is currently being considered by a number of ground, balloon and space based experiments, included the PLANCK experiment. The decomposition of the CMB polarization into $E$- and $B$-modes requires a full sky data coverage and, as such, is limited by the foreground contaminations. The latter introduce a mixing of the $E$ polarization into $B$ with the corresponding cosmic variance limitation. PLANCK’s expected sensitivity is about $T/S = 0.05$ corresponding to a minimum testable value of $H \simeq 6.7 \times 10^{13}\text{ GeV}$ [37]. A detection of the tensor mode would therefore imply that the value of the Hubble rate during inflation is larger than about $10^{13}\text{ GeV}$. As we will see, this will have important implications for the the ideas discussed in this paper.

The paper is organized as follows. In Section 2 we describe some mechanisms by which the SM Higgs could generate the cosmological perturbation. Being ignorant about the exact mechanism, we try to be as generic as possible. In Section 3 we present our results and draw our conclusions.
2 Curvature perturbation from the SM Higgs

Let us see in more detail how the curvature perturbation may be originated from the fluctuation of the Higgs field. A convenient and rather model-independent way to characterize the curvature perturbation generated during or after a de Sitter stage is by the $\delta N$ formalism [12,38,39]. The comoving curvature perturbation $\zeta$ on a uniform energy density hypersurface at time $t_f$ is, on sufficiently large scales, equal to the perturbation in the time integral of the local expansion from an initial flat hypersurface ($t = t_*$) to the final uniform energy density hypersurface. On sufficiently large scales, the local expansion can be approximated quite well by the expansion of the unperturbed Friedmann universe. Hence the curvature perturbation at time $t_f$ can be expressed in terms of the values of the relevant light scalar fields (the inflaton field, the SM Higgs, etc.) $\sigma^I(t_*, \vec{x})$ at $t_*$

$$\zeta(t_f, \vec{x}) = N_I \sigma^I + \frac{1}{2} N_{IJ} \sigma^I \sigma^J + \cdots \quad (I = 1, \ldots, M),$$

(1)

where $N_I$ and $N_{IJ}$ are the first and second derivative, respectively, of the number of e-folds

$$N(t_f, t_*, \vec{x}) = \int_{t_*}^{t_f} dt H(t, \vec{x})$$

(2)

with respect to the field $\sigma^I$. From the expansion (1) one can read off the the two-point correlator of the comoving curvature perturbation in momentum space

$$P_\zeta(k_1) = N_I N_J P_{k_1}^{IJ},$$

$$\langle \sigma^I_{k_1} \sigma^J_{k_2} \rangle = (2\pi)^3 \delta(k_1 + k_2) P_{k_1}^{IJ} = (2\pi)^3 \delta(k_1 + k_2) \delta^{IJ} \left( \frac{H}{2\pi} \right)^2.$$  

(3)

The last passage, stating that perturbations are not cross-correlated, holds if the cosmological perturbation is sourced by light scalar fields other than the inflaton as a consequence of the spatial conformal symmetry enjoyed by the de Sitter geometry [40]. The subscript $k_1$ is there to remind that perturbations have to be evaluated at horizon-crossing, when $k_1 = a H$, being $a$ the scale factor. From now on we restrict ourselves to the case of $M = 2$ and we identify one of the two fields with the adiabatic inflaton mode $\phi$ and the other with the SM Higgs. We also assume here that the contribution from the inflaton is subleading, $N_\phi \ll N_h$.

A specific example where the primordial density perturbations may be produced just after the end of inflation is the modulated decay scenario when the decay rate of the inflaton is a function of the SM Higgs field [30], that is $\Gamma = \Gamma(h)$. If we approximate the inflaton reheating by a sudden decay, we may find an analytic estimate of the density perturbation. In the case of modulated reheating, the decay occurs on a spatial hypersurface with variable local decay rate and hence local Hubble rate $H = \Gamma(h)$. Before the inflaton decay, the oscillating inflaton
field has a pressureless equation of state and there is no density perturbation. The perturbed expansion reads
\[ \delta N_d = -\frac{1}{3} \ln \left( \frac{\rho_\phi}{\rho_\phi} \right). \] (4)
Immediately after the decay we have radiation and hence the curvature perturbation reads
\[ \zeta = \delta N_d + \frac{1}{4} \ln \left( \frac{\rho_\phi}{\rho_\phi} \right). \] (5)
Eliminating \( \delta N_d \) and using the local Friedmann equation \( \rho \sim H^2 \), to determine the local density in terms of the local decay rate \( \Gamma = \Gamma(h) \), we have at the linear order
\[ \zeta = -\frac{1}{6} \frac{\delta \Gamma}{\Gamma} = -\frac{1}{6} \frac{\text{d} \ln \Gamma}{\text{d} \ln h} \frac{\delta h}{h} = -\beta_h \frac{\delta h}{h}. \] (6)
A maybe more intuitive way of understanding the expression (6) is to remember that, if the inflaton decay is perturbative, the final reheating temperature \( T_r \) scales like \( (M_{\text{Pl}} \Gamma)^{1/2} \) and therefore large scale spatial variations of the decay rate will induce a temperature anisotropy, \( \delta T_r/T_r \sim \delta \Gamma/\Gamma \). The corresponding power spectrum of the curvature perturbation is given by
\[ P_\zeta = \beta_h^2 \left( \frac{H}{2\pi \bar{h}} \right)^2. \] (7)
Another possibility is that that the dominant component of the curvature perturbation is generated at the transition between inflation and the post-inflationary phase [32]. Let us suppose that slow-roll inflation suddenly gives way to radiation domination through a waterfall transition in hybrid inflation and that the value of the inflaton \( \phi_e \) at which this happens depends on the SM Higgs field, \( \phi_e = \phi_e(h) \). If we assume again that the contribution to the curvature perturbation from the inflaton field is subdominant, we get
\[ \zeta = N'_e \delta \phi_e = N'_e \frac{\text{d} \phi_e}{\text{d} \ln h} \frac{\delta h}{h} = \frac{\text{d} N_e}{\text{d} \ln h} \frac{\delta h}{h}. \] (8)
and the power spectrum is given by Eq. (7) with \( \beta_h = \text{d} N_e/\text{d} \ln \bar{h} \). Given our ignorance about the parameter \( \beta_h \), in the following we will treat it as a free parameter. Specific forms of the function \( \beta_h \) may give rise to other interesting and observable properties, such as non-Gaussianity in the perturbation [41]. We will come back to these issues in the future.

In all the considerations made so far we have assumed that the Higgs field is quantum-mechanically excited during the de Sitter stage and acts like an isocurvature mode. This imposes various conditions:

1. The Higgs field does not contribute significantly to the energy density during inflation, \( V(h) \ll H^2 M_{\text{Pl}}^2 \).
2. The SM Higgs field is light enough during inflation, that is the second derivative of the SM Higgs potential \( V(h) \) is smaller than the square of the Hubble rate, \( |d^2V(h)/dh^2| \ll H^2 \). This condition has to hold for a number of e-folds large enough to create a sufficiently homogenous and isotropic region encompassing the comoving volume of the entire presently observable universe (we take 60 as a fiducial number) and also implies that the spectral index \( n_\zeta \) is sufficiently close to unity

\[
 n_\zeta - 1 = \frac{d\ln P_\zeta}{d\ln k} = \frac{\ln H_k^2}{\ln k} + \frac{2}{3H^2} \frac{d^2V(h)}{dh^2} = -2\epsilon + \frac{2}{3H^2} \frac{d^2V(h)}{dh^2}. \tag{9}
\]

We take \(|d^2V(h)/dh^2|/(3H^2) < 10^{-2}\) as a sufficient starting condition. This also automatically implies that \(|d\ln h/d\ln a| \ll 1\).

3 Results and Conclusions

We use 2-loop renormalization group (RG) equations for all SM couplings (gauge, Higgs quartic and top-yukawa couplings), and the pole mass matching scheme for the Higgs and top masses, as given in the Appendix of Ref. [22]. The numerical solution of these equations allows to obtain the RG-improved effective potential for the SM Higgs [42], as a function of input parameters, such as the Higgs and top masses. For the top mass we considered \(m_t = 173.1 \pm 0.7\) GeV (as in Ref. [24] by combining Tevatron [43] and LHC results), while for the QCD gauge coupling \(\alpha_s(M_Z) = 0.1184 \pm 0.0007\) [44]. In all simulations we have set \(\alpha_s\) to its central value; the size of the effect of the variation of \(\alpha_s\) within 1σ is comparable with the higher-order corrections (such as 3 loops) we are neglecting.

The Higgs field \(h\) is initially placed at \(h_{\text{initial}}\) where the second derivative of the potential is such that \(|d^2V(h)/dh^2|/(3H^2) = 10^{-2}\). Then, the field rolls down the potential according to the equation

\[
 \ddot{h} + 3H\dot{h} + V'(h) = 0, \tag{10}
\]

where the dot refers to derivative with respect to \(t = H^{-1}\ln a\). We have conventionally set the scale factor equal to 1 at the initial point \(h_{\text{initial}}\).

In Figure 1 we show an example of a situation where the conditions 1. and 2. in the previous section are met and the mechanism works. In fact, the field is slowly rolling down the potential, keeping the second derivative of the potential small (in units of \(H^2\)), over a wide range of e-folds. So, in this case it is possible to generate nearly scale-independent isocurvature perturbations of the SM Higgs and, through one of the mechanisms mentioned above, convert them into the observed amount of curvature perturbations \(P_\zeta\).

Next, we repeat the analysis by scanning over the Higgs and top masses and looking for what values of \(H\) the same situation arises. Of course, an upper limit on the values of \(H\) is set
by the instability scale $\Lambda_{\text{inst}}$ at which the Higgs quartic coupling runs negative. For lower values of $H$, down to $\sim 1$ TeV, we verified that there always exist solutions satisfying conditions 1 and 2 of the previous section, and generating enough curvature perturbations. The analysis of the region $H \lesssim 1$ TeV is more delicate as the Higgs field runs very close to its minimum, but we do not pursue this here (we are more interested in the regime where $H$ is relatively large so that tensor modes could be seen by PLANCK).

In Figure 2 we show the resulting upper limits on $H$, corresponding to the instability scale $\Lambda_{\text{inst}}$, for a range of $m_h$ around the observed value. The different lines refer to $m_t$ at the central value or within $1\sigma$ and $2\sigma$ variations. The mechanism proposed in this paper works for all values of $H$ below the curves. The current exclusion limit on $H$ from CMB data is $H \lesssim 1.5 \times 10^{14}$ GeV [36], which is dark gray band in Fig. 2). As already mentioned, PLANCK data will possibly be able to exclude values of the Hubble rate down to $H \simeq 6.7 \times 10^{13}$ GeV. This limit is shown as a lighter gray band in Fig. 2. Let us reiterate that the values of $\beta_h$, see Eq. (7), is not fixed, but is appropriately computed at every point to insure that the cosmological perturbation is correctly normalized, $P_{\zeta}^{1/2} = 4.8 \times 10^{-5}$. In other words, Figure 2 provides all the possible values of $H$ where the perturbation may be generated by the SM Higgs.

The information in Figure 2 can be read in two ways. For given $m_h$ and $m_t$, there is a maximum $H$ for which the SM Higgs can generate cosmological perturbations, and which may
Figure 2: Upper limits on $H$ as a function of $m_h$. Solid black line refers to central value of $m_t$, while dotted blue and dashed red lines correspond to 1$\sigma$ and 2$\sigma$ variations of $m_t$, respectively.

or may not be in the range to be detected in the near future. On the other hand, for a given value of $H$, the hypothesis that the SM Higgs generate perturbations establishes a correlation between $m_h$ and $m_t$ which can be tested by particle physics measurements.

For example, if the Higgs mass value is confirmed to be $m_h = 125.5$ GeV and $m_t$ and $\alpha_s$ are at their central values, our mechanism predicts $H \lesssim 10^{10}$ GeV and therefore tensor perturbations too small to be detected in the near future. On the other hand, if tensor perturbations will be detected with a given $H$, the mechanism we have proposed makes a prediction for a relation between the Higgs and top masses, and we find

$$(m_h)^{\text{B-mode}} \simeq 128.0 \text{ GeV} + 1.3 \left(\frac{m_t - 173.1 \text{ GeV}}{0.7 \text{ GeV}}\right) \text{ GeV} + 0.9 \left(\frac{H}{10^{15} \text{ GeV}}\right) \text{ GeV} \pm \delta_{\text{th}}, \quad (11)$$

where $\delta_{\text{th}} \sim 2$ GeV is a residual theoretical uncertainty from higher order corrections, we have neglected. The formula (11) is valid for $124 \text{ GeV} \lesssim m_h \lesssim 127 \text{ GeV}$ and $6.7 \times 10^{13} \text{ GeV} \lesssim H \lesssim 1.5 \times 10^{14} \text{ GeV}$. The information coming from a more accurate experimental determinations of the SM parameters would then allow to either support or rule out our model. Notice also that the result in Eq. (11) is crude, and just serves as an illustration. A more careful calculation (also including 3-loop $\beta$-functions and $\alpha_s$ variation) would be needed in order to extract a detailed and more complete prediction.

Since the purpose of this paper is to show whether it is possible to generate cosmological perturbations with the SM Higgs, the level of accuracy we adopted is enough to conclude that the answer is robustly yes. A more accurate determination of the Higgs effective potential,
using 3-loop $\beta$-functions for gauge, Higgs quartic and top-yukawa couplings \cite{45} and 2-loop pole mass matching conditions (as recently presented in Refs. \cite{24,25}), would be desirable, but is beyond the scope of the present brief communication.

In conclusion, we have pointed out the possibility that the SM Higgs might be responsible for the inhomogeneities we observe in our universe, both in the CMB and in the LSS, if there was an early stage of accelerated expansion driven by some vacuum energy and whose size we agnostically parametrized with the Hubble rate $H$. Essentially, the Higgs potential may be flat enough to generate nearly scale-independent isocurvature perturbations which are subsequently converted into the observed adiabatic mode. With the recent ATLAS and CMS results for the Higgs mass, the Higgs can generate the cosmological perturbation in a wide range of the Hubble rate during inflation, $H = (10^{10} - 10^{14})$ GeV, depending on the values of the SM parameters. On the other side, if the forthcoming data from the Planck satellite will present some hints of a $B$-mode in the CMB polarization originated from tensor modes, this will identify a well-defined range of the Higgs mass. We have therefore established a very interesting correlation between collider and cosmological measurements, which makes the mechanism predictive and falsifiable.

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