Can Standard Model Higgs Seed the Formation of Structures in Our Universe?

Ki-Young Choi\textsuperscript{1}\textsuperscript{*} and Qing-Guo Huang\textsuperscript{2}\textsuperscript{†}

\textsuperscript{1}Asia Pacific Center for Theoretical Physics, Pohang, Gyeongbuk 790-784, Republic of Korea and Department of Physics, POSTECH, Pohang, Gyeongbuk 790-784, Republic of Korea and

\textsuperscript{2}State Key Laboratory of Theoretical Physics, Institute of Theoretical Physics, Chinese Academy of Science, Beijing 100190, People’s Republic of China

Abstract

We study the Standard Model Higgs field as a source for the primordial curvature perturbation, particularly in the curvaton and modulated reheating scenario. We conclude that the Higgs cannot play as a curvaton due to the small energy density when it decays, however the modulated reheating by Higgs can be a viable scenario. In the latter case, the non-Gaussianity is inevitably generated and strongly constrains the type of potential of inflaton field and Higgs-dependent interaction term. For the quadratic potential of the inflaton field with decay rate which non-linearly depends on the Higgs vacuum expectation value, the contribution of Higgs field to the primordial curvature perturbation must be less than 8%.

PACS numbers: 98.80.-k,98.80.Cq,98.80.Es

\textsuperscript{*}kiyoung.choi@apctp.org
\textsuperscript{†}huanggg@itp.ac.cn
I. INTRODUCTION

Cosmic inflation solves various problems in the standard Big Bang cosmology [1] and simultaneously provides the seed of large scale structure in our Universe from the quantum fluctuation of a light scalar field [2].

Recently ATLAS and CMS collaborations at CERN reported a discovery of Higgs-like particle with the mass of 125 GeV [3, 4]. As a unique scalar field in the Standard Model (SM), Higgs can be considered as a source field for the inflation. However it is known that its large self-interaction predicts very large density perturbation, which is easily ruled out by observation [5, 6].

It was recently demonstrated [7, 8] that the SM itself can give rise to inflation, provided non-minimal coupling of the Higgs field with gravity. However this model is plagued by several issues such as unitarity problem [9], and the stability of the potential up to the near Planck scale [6]. The sensitivity of the Higgs inflation scenario to the details on the UV completion of the theory was studied recently in [10].

Even though the canonical Higgs field in the Standard Model is not a good candidate for inflation, it may contribute to the primordial curvature perturbation to seed the structure formation and anisotropies in cosmic microwave background radiation. The Higgs perturbation is generated during inflation on super horizon scales at horizon exit if its mass is smaller than the Hubble scale. During inflation this remains the isocurvature perturbation, however it can be converted to the adiabatic one after or at the end of inflation. In this case Higgs can contribute to the primordial curvature perturbation.

In the Ref. [11], the authors studied the possibility that the SM Higgs might be responsible for the inhomogeneties we observe in our universe in the alternative scenarios such as modulated decay [12] and inhomogeneous end of inflation [13, 14]. In particular they discussed about the implications for the detection of primordial tensor perturbations in the future observation of $B$-mode of CMB polarization.

In this paper we study the idea of the Higgs field as a source for the primordial curvature perturbation, however we focus on the observables from the non-linear terms in the curvature perturbations, the non-Gaussianity. Wilkinson Microwave Anisotropy Probe (WMAP) seven-year data constrains the local type nonlinear parameter, $-10 < f_{NL} < 74$ at the 95% confidence level [15].
We find that Higgs as a curvaton is difficult since the energy density of the Higgs condensate is too small at the time of energy transfer of Higgs to the radiation. In the modulated reheating scenario, Higgs can contribute to the primordial curvature perturbation as shown in [11]. However we find that the Higgs-dependent decay can generate large non-Gaussianity. For the chaotic inflation with quadratic potential, too large non-Gaussianity is generated to be compatible with the current observation due to the non-linearly Higgs-dependent decay. In this case, we find that the Higgs contribution to the primordial curvature perturbation should be less than 8%.

The paper is organized as follows. In Section II we summarize the properties of Higgs field during inflation. In Section III we study the possibility of Higgs as curvaton. In Section IV we consider the modulated reheating by Higgs. In Section V we summarize our results.

II. HIGGS FIELD DURING INFLATION

We consider the SM Higgs potential in the form,
\begin{equation}
V(h) = -\frac{1}{2}m^2 h^2 + \frac{\lambda}{4} h^4,
\end{equation}
where the mass of Higgs in the minimum is \( m_h = \sqrt{2}m \simeq 125 \text{ GeV} \) and the self coupling \( \lambda \equiv \lambda(\mu) \) has a logarithmic dependence on the energy scale. When we use the central values for the top quark mass and strong coupling constant, the Higgs potential develops an instability around \( 10^{11} \text{ GeV} \), with a lifetime much longer than the age of the Universe. However taking into account theoretical and experimental errors, stability up to the Planck scale cannot be excluded [16–18]. Therefore in this paper it is legitimate to assume that the Higgs potential is stable up to the energy scale of our interest below Planck scale.

Since the canonical Higgs field is not suitable for inflation, we assume that there is another scalar field which drives inflation at very high energy scale. The scale of the inflation is constrained by the non-observation of the tensor spectrum. The present bound on the tensor-to-scalar ratio gives upper bound on the Hubble parameter during inflation [15],

\begin{equation}
H \lesssim 3 \times 10^{14} \text{ GeV}.
\end{equation}

At this energy scale of inflation which is much larger than the electroweak scale, it is rea-
sonable to simplify Higgs potential Eq. (1) to be quartic as

\[ V(h) \simeq \frac{\lambda}{4} h^4, \]  

where the size of the Higgs self coupling is \( \lambda(\mu) \simeq \mathcal{O}(10^{-2}) \) \([16, 18]\). For simplicity we take \( \lambda \simeq 0.01 \) and constant at the scale of our study.

For the generation of the quantum fluctuations of the Higgs field, we require that the effective mass of the Higgs,

\[ m_{\text{eff}}^2 = V_{hh} = 3\lambda h^2, \]  

is smaller than the Hubble parameter during inflation. This condition sets the upper bound on the Higgs VEV

\[ h \ll \frac{H}{\sqrt{3\lambda}} \lesssim 2 \times 10^{15} \text{GeV}, \]  

where we used Eq. (2) and \( \lambda = 0.01 \). Because Higgs is quite light \(^1\), it almost does not move from its initial position and stays there during inflation. At this stage, one can easily check that the Higgs energy density is subdominant compared to the inflaton energy density by the order of less than \( H^2/M_P^2 \) as long as Eq. (5) is satisfied.

On the other hand, the Higgs fluctuations is generated during inflation with the amplitude

\[ \delta h = \frac{H}{2\pi}. \]  

Combining with Eq. (5), we lead to the range of the Higgs perturbation to Higgs VEV

\[ 0.03 \simeq \frac{\sqrt{3\lambda}}{2\pi} \lesssim \frac{\delta h}{h} \lesssim 1, \]

which is independent on the precise value of Hubble parameter. Here the upper bound comes by considering that the VEV of Higgs should be larger than the amplitude of its quantum fluctuations.

\(^1\) The Higgs is plagued by the quadratic divergence of its mass, but in our discussion we assume that the hierarchy problem is solved in some way and we don’t address this problem.
III. HIGGS FIELD AS CURVATON

A light scalar field which is subdominant during inflation can decay very late after inflation, when its vacuum energy is transfer to the radiation. At this time the isocurvature perturbation of curvaton field can be converted to the adiabatic one and contribute to the primordial curvature perturbation. This is known as curvaton scenario \[19\text{–}23\]. There have been many studies on this with different types of potentials of curvaton \[24\] including the inflaton contribution \[25\] to study the power spectrum and non-Gaussianities \[26\]. Since Higgs is subdominant during inflation and decay late it can be a natural candidate for curvaton. In this section we study the possibility of the SM Higgs field as curvaton.

Inflation is assumed to be driven by another scalar field, inflaton. During inflation the energy density of Higgs is subdominant and it acquires perturbation as explained in section II. After inflation, the inflaton field oscillates around its local minimum and its potential energy is converted into radiations to initiate the radiation-dominated Universe. During this time, Higgs field starts oscillation when the effective mass becomes bigger than the Hubble expansion, namely when

\[
H^2_{\text{osc}} = m^2_{\text{eff}} = 3\lambda h^2.
\]  

(8)

Since the potential of Higgs is quartic, the energy density of an oscillating Higgs decreases as \(a^{-4}\) with scale factor \(a\) which is similar to radiations. Therefore during the epoch of the oscillation of Higgs, the ratio of Higgs energy density to that of the radiation-dominated background continues to be constant. Actually, after some oscillations of Higgs, the energy of Higgs condensate is transferred to the radiation by perturbative and non-perturbative processes. Since the coupling of Higgs to other fields are strong the energy transfer is completed soon after the start of oscillation \[27\text{,} 28\]. Therefore the Higgs density parameter \(\Omega_h\) at the time of Higgs decay is approximatetely the same as the onset of oscillation of Higgs, which is given by

\[
\Omega_h = \frac{\rho_h}{\rho_r} = \left(\frac{\rho_h}{\rho_r}\right)_{\text{osc}} = \frac{\lambda/4 h^4_{\text{osc}}}{3M^2_P H^2_{\text{osc}}} = \frac{h^2}{36M^2_P} \ll \frac{H^2}{108\lambda M^2_P} \lesssim 1.5 \times 10^{-8}.
\]  

(9)

Here we used \(h_{\text{osc}} = h\), the higgs value during inflation and the inequality comes using Eq. (2) and Eq. (5). The density perturbation generated after the decay of Higgs depends
on the energy density of the Higgs when it decay, and it is given by
\[
\frac{\delta \rho}{\rho} \sim \frac{\rho_h \delta \rho_h}{\rho} \sim \frac{\rho_h \delta h}{h} \sim \Omega_h \frac{\delta h}{h}.
\] (10)

The Higgs perturbation $\delta h/h$ is already bounded to be less than $O(1)$ in Eq. (7). Therefore we can easily see that, due to the small energy density of Higgs in Eq. (9), it cannot produce enough density perturbation which is necessary for observation today. Thus we conclude that Higgs cannot play as curvaton. In the Ref. [29], the authors also argued that Higgs is not viable as curvaton in the context of generalized G inflation model.

IV. MODULATED REHEATING BY HIGGS

After inflation the energy density in the inflaton field must be transferred into radiation. In the simplest case of single field inflation model, reheating process does not affect the primordial curvature perturbation on scales which are observable today, because these scales were much larger than the horizon at the time of reheating. However when there is a subdominant light scalar field which modulates the efficiency of the reheating the situation changes [12, 30, 31]. The quasi-scale invariant perturbations in this light field, which during inflation are an isocurvature perturbation, can be converted into the primordial curvature perturbation during this process [32]. In this section, we examine the possibility of Higgs as a modulating particle in the reheating after inflation.

We assume that the inflaton field, $\phi$, drives inflation and its subsequent decay during oscillation transfers its vacuum energy to the background radiation. We consider the Lagrangian
\[
\mathcal{L} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) + \mathcal{L}_{\text{int}},
\] (11)
where $V(\phi)$ is the potential which is responsible for the inflation and $\mathcal{L}_{\text{int}}$ is the interaction responsible for the inflaton decay, which depends on the Higgs VEV, $h$. For simplicity we will take a polynomial potential for inflaton, $V(\phi) \propto \phi^{2\alpha}$. For the Higgs dependent interactions we consider [33]
\[
\mathcal{L}_{\text{int}} \supset - \sum_a y_a(h) \phi \bar{\psi}_a \psi_a - \sum_a M(h) \phi \chi_a^2 - \sum_a g_a(h) \phi^2 \chi_a^2,
\] (12)
where $\chi_a$ and $\psi_a$ are scalar and fermion fields which contribute to the radiation in the early Universe. The coupling constants $y_a(h), M_a(h)$, and $g_a(h)$ are functions of the Higgs field $h$. 
The total decay rate of the inflaton can be composed of Higgs independent and dependent part,
\[ \Gamma(h) = \Gamma^I + \Gamma^D(h). \]  
(13)

Using Eq. (12), Higgs dependent decay rate \( \Gamma^D(h) \) has the form of \[ \[33\]
\[ \Gamma^D(h) = \sum a_n A_n y_n^2(h) \frac{m_{\phi}^{\text{eff}}}{8\pi} + \sum B_n M_n(h) \frac{m_{\phi}^{\text{eff}}}{8\pi m_{\phi}^{\text{eff}}} + \sum C_n \frac{h_n^2(h)}{8\pi (m_{\phi}^{\text{eff}})^3} \rho_{\phi}, \]  
(14)

with \( m_{\phi}^{\text{eff}} \) the effective mass of the inflaton field defined by \( (m_{\phi}^{\text{eff}})^2 = \partial^2 V / \partial^2 \phi \) and \( A_n, B_n, \) and \( C_n \) are numerical coefficients of the order of \( \mathcal{O}(1 - 1000) \) \[33\].

After the reheating process, the primordial curvature perturbation \( \zeta \) has two contributions from the inflaton and the Higgs field. Due to the slow-rolling the inflaton field does not contribute to the curvature perturbation at non-linear orders. In contrast the modulating field can easily generate the non-linear contribution to the \( \zeta \). In this mixed modulating scenario, the \( \zeta \) can be written as \[31, 33\]
\[ \zeta = \frac{1}{M_P^2 V_{\phi}} \delta \phi_* + Q_h \delta h_* + \frac{1}{2} Q_{hh} \delta h_*^2 + \frac{1}{6} Q_{hhh} \delta h_*^3 + \cdots, \]  
(15)

where \( Q \) is a function of \( \Gamma(h)/H_c \) calculated at a time \( t_c \) which is after several oscillations of the inflaton but well before the time of decay of inflaton. A quantity with subscript \( * \) is evaluated when the corresponding scale crosses the Hubble horizon during inflation. For \( \Gamma/H_c \ll 1 \), \( Q \) can be well approximated by \[31, 33\]
\[ Q = a_0 \log \left( \frac{\Gamma}{H_c} \right), \]  
(16)

where \( a_0 \) has different value and sign for different inflaton potential and interaction for the dominant decay mode. The value of \( a_0 \) for some of the examples are given in the Table I of Ref. [33]. With this form we find the derivatives
\[ Q_h = a_0 \frac{\Gamma_h}{\Gamma}, \]
\[ Q_{hh} = a_0 \left( \frac{\Gamma_{hh}}{\Gamma} - \frac{\Gamma_h^2}{\Gamma^2} \right), \]  
(17)
\[ Q_{hhh} = a_0 \left( \frac{\Gamma_{hhh}}{\Gamma} - 3 \frac{\Gamma_h \Gamma_{hh}}{\Gamma^2} + 2 \frac{\Gamma_h^3}{\Gamma^3} \right). \]

Once \( \zeta \) in Eq. (15) is given, we can easily calculate the power spectrum of curvature perturbation
\[ P_\zeta = P_{\zeta\phi} + P_{\zeta h} = \frac{1}{2M_P^2 \epsilon_*} \left( \frac{H_*}{2\pi} \right)^2 (1 + \tilde{r}). \]  
(18)
The power spectrum has contributions from each field and we define the ratio by $\tilde{r}$ as

$$\tilde{r} \equiv \frac{P_{\zeta_h}}{P_{\zeta}} = 2M_P^2\epsilon_*Q_h^2,$$  \hspace{1cm} (19)

where $\epsilon$ is the slow-roll parameter

$$\epsilon \equiv \frac{M_P^2}{2} \left( \frac{V_{\phi}}{V} \right)^2.$$  \hspace{1cm} (20)

The spectral index of power spectrum of curvature perturbation is also obtained

$$n_s \equiv 1 + \frac{d \ln P_{\zeta}}{d \ln k} = (1 - \beta)n_{\zeta} + \beta n_{\zeta h},$$  \hspace{1cm} (21)

where $\beta \equiv \tilde{r}/(1 + \tilde{r}) \in [0, 1]$. Here we used the notations

$$n_{\zeta} = 1 - 6\epsilon + 2\eta_{\phi\phi},$$  \hspace{1cm} (22)

$$n_{\zeta h} = 1 - 2\epsilon + 2\eta_{hh},$$  \hspace{1cm} (23)

and

$$\eta_{\phi\phi} \equiv M_P^2 \frac{V_{\phi\phi}}{V}, \quad \eta_{hh} \equiv \frac{V_{hh}}{3H^2}. \hspace{1cm} (24)$$

The tensor perturbation depends only on the total energy density during inflation, and thus the amplitude of its power spectrum takes the form

$$P_T = \frac{H_*/M_P^2}{\pi^2/2},$$  \hspace{1cm} (25)

whose tilt is given by

$$n_T \equiv \frac{d \ln P_T}{d \ln k} = -2\epsilon_*.$$  \hspace{1cm} (26)

Usually we introduce a new parameter, so-called tensor-to-scalar ratio $r$, to measure the size of gravitational wave perturbation,

$$r \equiv \frac{P_T}{P_{\zeta}} = (1 - \beta)16\epsilon_*.$$  \hspace{1cm} (27)

And thus we obtain the consistency relation scenario between $n_T$ and $r$ for inflaton-modulated reheating,

$$n_T = \frac{1}{1 - \beta \frac{r}{8}}.$$  \hspace{1cm} (28)

In the limit of $\beta \to 0$, the consistency relation becomes the prediction of the usual single-field slow-roll inflation model.
Without loss of generality, we can assume that $\Gamma^D(h)$ is proportional to polynomial of the Higgs field, $h^n$, with $n = 1, 2, 3, \cdots$. For each case of $n$, $Q_h$ in Eq. (17) takes the form

$$Q_h = a_0 B_h \frac{n}{h},$$

(29)

where $B_h \equiv \Gamma^D/\Gamma$, the fraction of the Higgs-dependent decay rate to the total decay rate of inflaton field. Using the observed power spectrum, $P_\zeta = 2.46 \times 10^{-9}$ \cite{15}, and Eq. (18), then $B_h$ has to satisfy the relation

$$B_h = 3.1 \times 10^{-4} \frac{\sqrt{\beta}}{n|a_0|} \left( \frac{h_s}{H_s} \right) \ll 1.7 \times 10^{-3} \frac{\sqrt{\beta}}{n|a_0|},$$

(30)

where the last inequality comes from Eq. (5). Considering that $\sqrt{\beta} \sim \mathcal{O}(10^{-1})$ and $a_0 \sim 0.1$, we conclude that

$$B_h \lesssim \mathcal{O}(10^{-3} \sim 10^{-2}).$$

(31)

It indicates that inflaton decay is dominated by the Higgs independent part.

During the modulated reheating process by Higgs, the non-linear term of $\zeta$ in Eq. (15) can generate large non-Gaussianity, which might be inconsistent with observational bound. From Eq. (15), we can calculate the non-linearity parameters $f_{NL}$, $\tau_{NL}$, and $g_{NL}$. Following \cite{33}, in the leading order of the slow-roll parameters, they are

$$f_{NL} = \frac{5 \beta^2}{6 a_0} \left( -1 + \frac{\Gamma_{hh}}{\Gamma_h^2} \right),$$

$$\tau_{NL} = \frac{\beta^3}{a_0^2} \left( 1 - \frac{\Gamma_{hh}}{\Gamma_h^2} \right)^2,$$

$$g_{NL} = \frac{25 \beta^3}{54 a_0^2} \left( 2 - 3 \frac{\Gamma_{hh}}{\Gamma_h^2} + \frac{\Gamma_{hh}}{\Gamma_h^2} \right).$$

(32)

Taking into account of the polynomial form of the Higgs-depenent decay rate, $\Gamma^D(h) \propto h^n$, $f_{NL}$ can be written by

$$f_{NL} \simeq \frac{-5 \beta^2}{6 a_0} + \frac{5 \beta^2}{6 a_0 B_h} \frac{n - 1}{n}.$$

(33)

When the linear term dominates the Higgs-dependent term ($n = 1$), only the first term survives and $f_{NL} \simeq \frac{-5 \beta^2}{6 a_0}$, which is smaller than $\mathcal{O}(1 \sim 10)$. This is consistent with observation. However when the non-linear term dominates ($n > 1$), $f_{NL}$ is mostly determined by the second term in Eq. (33) and can be large due to the small $B_h$. For this case with large
If \( \beta \ll 1 \), \( \tau_{NL} \) are significantly enhanced compared to \( f_{NL}^2 \).

For \( n = 1 \), we have \( a_0 = -1/6 \) if the inflaton potential is quadratic, \( V(\phi) \propto \phi^2 \). The non-linearity parameter \( f_{NL} = 5\beta^2 \leq 5 \) and then the current bound on \( f_{NL} \) does not give any constraint on \( \beta \); if \( V(\phi) \propto \phi^6 \) we have \( a_0 = 1/6, 1/30, 1/18 \) corresponding to the dominant interaction for inflaton decay \( -y\phi\bar{\psi}\psi, -M\phi\chi^2, -\lambda\phi^2\chi^2 \) respectively, and the bound on \( f_{NL} \) implies

\[
\beta < \min(\sqrt{12a_0}, 1).
\]

We see that the non-Gaussianity does not constrain the model and the Higgs can significantly contribute to the curvature perturbation in the case of \( n = 1 \).

However when the non-linearly dependent term dominates the result changes. For \( n > 1 \), with the quadratic type of inflaton potential, \( V(\phi) \propto \phi^2 \), the present bound on the nonlinear parameter \( f_{NL} \) gives strong constraint. We obtain the upper bound \( \beta^2 < \frac{12n}{n-1} |a_0|B_h \) for negative \( a_0 \). Combining with Eq. (30), this gives

\[
\beta \lesssim \frac{0.075}{(n-1)^{2/3}}.
\]

Since \( \beta \simeq \tilde{r} \) for small \( \beta \), the constraint on \( \beta \) translates into the bound on the contribution of the Higgs field to the power spectrum. From Eq. (37), we conclude that the Higgs can contribute only 8% to the power spectrum of the primordial curvature perturbation and this is very subdominant.

V. DISCUSSION

We have examined the viability of the Standard Model Higgs as a dominant source of the primordial curvature perturbation. We find that Higgs as a curvaton is difficult since the energy density of the Higgs condensate is too small at the time of energy transfer of Higgs
to the radiation. In the modulated reheating scenario, Higgs can dominantly contribute to
the perturbation when the Higgs-dependent decay rate is linearly proportional to the Higgs
VEV. In this case the primordial non-Gaussianity generated by Higgs must be small. On
the other hand, if the Higgs-dependent decay rate depends non-linearly on the Higgs VEV,
too large non-Gaussianity is predicted. The current bound on the non-linearity parameter
shows that the contribution of the Higgs to the total curvature perturbation must be smaller
than 8% for the quadratic potential of inflaton field.

Acknowledgments

K.-Y.C was supported by Basic Science Research Program through the National Research
Foundation of Korea (NRF) funded by the Ministry of Education, Science and Technology
(No. 2011-0011083). K.-Y.C acknowledges the Max Planck Society (MPG), the Korea
Ministry of Education, Science and Technology (MEST), Gyeongsangbuk-Do and Pohang
City for the support of the Independent Junior Research Group at the Asia Pacific Center
for Theoretical Physics (APCTP). K.-Y.C would like to thank the ITP of CAS for warm
hospitality during his stay where this work was initiated. QGH is supported by the project
of Knowledge Innovation Program of Chinese Academy of Science and a grant from NSFC
(grant NO. 10975167).

[1] A. A. Starobinsky, JETP Lett. 30 682 (1979) [Pisma Zh. Eksp. Teor. Fiz. 30 719 (1979)];
K. Sato, Mon. Not. Roy. Astron. Soc. 195, 467 (1981);
A. H. Guth, Phys. Rev. D 23, 347 (1981);
A. D. Linde, Phys. Lett. B 108 389 (1982);
A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. 48 1220 (1982).
[2] S. W. Hawking, Phys. Lett. B 115, 295 (1982);
A. A. Starobinsky, Phys. Lett. B 117, 175 (1982);
A. H. Guth and S. Y. Pi, Phys. Rev. Lett. 49, 1110 (1982).
[3] G. Aad et al. [ATLAS Collaboration]. [arXiv:1207.7214 [hep-ex]].
[4] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B [arXiv:1207.7235 [hep-ex]].
[5] A. D. Linde, Chur, Switzerland: Harwood (1990) 362 p. (Contemporary concepts in physics, 5) [hep-th/0503203], D. H. Lyth and A. Riotto, Phys. Rept. 314 (1999) 1 [hep-ph/9807278], A. D. Linde, Lect. Notes Phys. 738 (2008) 1 [arXiv:0705.0164 [hep-th]].

[6] G. Isidori, V. S. Rychkov, A. Strumia and N. Tetradis, Phys. Rev. D 77 (2008) 025034 [arXiv:0712.0242 [hep-ph]].

[7] J. L. Cervantes-Cota and H. Dehnen, Nucl. Phys. B 442 (1995) 391 [astro-ph/9505069].

[8] F. L. Bezrukov and M. Shaposhnikov, Phys. Lett. B 659 (2008) 703 [arXiv:0710.3755 [hep-th]].

[9] C. P. Burgess, H. M. Lee and M. Trott, JHEP 0909 (2009) 103 [arXiv:0902.4465 [hep-ph]], J. L. F. Barbon and J. R. Espinosa, Phys. Rev. D 79 (2009) 081302 [arXiv:0903.0355 [hep-ph]], C. P. Burgess, H. M. Lee and M. Trott, JHEP 1007 (2010) 007 [arXiv:1002.2730 [hep-ph]].

[10] F. Bezrukov, A. Magnin, M. Shaposhnikov and S. Sibiryakov, JHEP 1101 (2011) 016 [arXiv:1008.5157 [hep-ph]].

[11] A. De Simone and A. Riotto, arXiv:1208.1344 [hep-ph].

[12] G. Dvali, A. Gruzinov and M. Zaldarriaga, Phys. Rev. D 69 023505 (2004).

[13] D. H. Lyth, JCAP 0511 (2005) 006. [astro-ph/0510443].

[14] D. H. Lyth and A. Riotto, Phys. Rev. Lett. 97 (2006) 121301 [astro-ph/0607326].

[15] E. Komatsu et al., Astrophys. J. Suppl. 192, 18 (2011).

[16] J. Elias-Miro, J. R. Espinosa, G. F. Giudice, G. Isidori, A. Riotto and A. Strumia, Phys. Lett. B 709 (2012) 222 [arXiv:1112.3022 [hep-ph]].

[17] F. Bezrukov, M. Y. .Kalmykov, B. A. Kniehl and M. Shaposhnikov, JHEP 1210 (2012) 140 [arXiv:1205.2893 [hep-ph]].

[18] G. Degrassi, S. Di Vita, J. Elias-Miro, J. R. Espinosa, G. F. Giudice, G. Isidori and A. Strumia, JHEP 1208 (2012) 098 [arXiv:1205.6497 [hep-ph]].

[19] S. Mollerach, Phys. Rev. D 42 313 (1990).

[20] A. D. Linde and V. F. Mukhanov, Phys. Rev. D 56 535 (1997).

[21] D. H. Lyth and D. Wands, Phys. Lett. B 524, 5 (2002).

[22] T. Moroi and T. Takahashi, Phys. Lett. B 522, 215 (2001) [Erratum-ibid. B 539, 303 (2002)].

[23] K. Enqvist and M. S. Sloth, Nucl. Phys. B 626, 395 (2002).

[24] K. Dimopoulos, G. Lazarides, D. Lyth and R. Ruiz de Austri, Phys. Rev. D 68 123515 (2003); K. Enqvist and S. Nurmi, JCAP 0510 013 (2005); Q. G. Huang, Phys. Lett. B 669, 260 (2008);
K. Enqvist and T. Takahashi, JCAP 0809, 012 (2008);
Q. G. Huang and Y. Wang, JCAP 0809 025 (2008);
Q. G. Huang, JCAP 0811, 005 (2008);
M. Kawasaki, K. Nakayama and F. Takahashi, JCAP 0901, 026 (2009);
P. Chingangbam and Q. G. Huang, JCAP 0904, 031 (2009);
K. Enqvist, S. Nurmi, G. Rigopoulos, O. Taanila and T. Takahashi, JCAP 0911, 003 (2009);
K. Enqvist and T. Takahashi, JCAP 0912, 001 (2009);
K. Enqvist, S. Nurmi, O. Taanila and T. Takahashi, JCAP 1004, 009 (2010);
Q. G. Huang, JCAP 1011, 026 (2010) [Erratum-ibid. 1102, E01 (2011)];
C. T. Byrnes, K. Enqvist and T. Takahashi, JCAP 1009, 026 (2010);
K. -Y. Choi and O. Seto, Phys. Rev. D 82 103519 (2010);
J. Fonseca and D. Wands, Phys. Rev. D 83, 064025 (2011);
C. T. Byrnes, K. Enqvist, S. Nurmi and T. Takahashi, JCAP 1111, 011 (2011);
M. Kawasaki, T. Kobayashi and F. Takahashi, Phys. Rev. D 84, 123506 (2011) [Phys. Rev. D 85, 029905 (2012)];
A. Mazumdar and J. Rocher, Phys. Rept. 497, 85 (2011) [arXiv:1001.0993 [hep-ph]].

[25] D. Langlois and F. Vernizzi, Phys. Rev. D 70, 063522 (2004);
G. Lazarides, R. R. de Austri and R. Trotta, Phys. Rev. D 70, 123527 (2004);
F. Ferrer, S. Rasanen and J. Valiviita, JCAP 0410, 010 (2004);
T. Moroi, T. Takahashi and Y. Toyoda, Phys. Rev. D 72 023502 (2005);
T. Moroi and T. Takahashi, Phys. Rev. D 72 023505 (2005);
K. Ichikawa, T. Suyama, T. Takahashi and M. Yamaguchi, Phys. Rev. D 78, 023513 (2008);
K. -Y. Choi and O. Seto, Phys. Rev. D 85 (2012) 123528.

[26] D. H. Lyth, C. Ungarelli and D. Wands, Phys. Rev. D 67, 023503 (2003), N. Bartolo, S. Matarrese and A. Riotto, Phys. Rev. D 69 043503 (2004), D. H. Lyth and Y. Rodriguez, Phys. Rev. Lett. 95 121302 (2005), M. Sasaki, J. Valiviita and D. Wands, Phys. Rev. D 74 103003 (2006), K. A. Malik and D. H. Lyth, JCAP 0609 008 (2006), K. Y. Choi and J. O. Gong, JCAP 0706 007 (2007).

[27] F. Bezrukov, D. Gorbunov and M. Shaposhnikov, JCAP 0906 (2009) 029 [arXiv:0812.3622 [hep-ph]].

[28] J. Garcia-Bellido, D. G. Figueroa and J. Rubio, Phys. Rev. D 79 (2009) 063531
[29] T. Kunimitsu and J. ’i. Yokoyama, arXiv:1208.2316 [hep-ph].

[30] L. Kofman, arXiv:astro-ph/0303614

[31] M. Zaldarriaga, Phys. Rev. D 69, 043508 (2004).

[32] L. Alabidi, K. Malik, C. T. Byrnes and K. -Y. Choi, JCAP 1011 (2010) 037 arXiv:1002.1700 [astro-ph.CO]]

[33] K. Ichikawa, T. Suyama, T. Takahashi, M. Yamaguchi, Phys. Rev. D78 063545 (2008).

[34] http://www.rssd.esa.int/index.php?project=PLANCK