Lowering the Self-Coupling of the Scalar Field in the Generalized Higgs Inflation

Kourosh Nozari\(^1\) • Somayeh Shafizadeh\(^2\) • Narges Rashidi\(^3\)

Abstract We study cosmological dynamics of a generalized Higgs inflation. By expanding the action up to the second and third order in the small perturbations, we study the primordial perturbation and its non-Gaussian distribution. We study the non-Gaussian feature in both the equilateral and orthogonal configurations. By adopting a quartic potential, we perform a numerical analysis on the model’s parameter space and compare the results with Planck2015 observational data. To obtain some observational constraint, we focus on the self-coupling and the non-minimal coupling parameters. We show that, in the presence of the non-minimal coupling and the Galileon-like interaction, the self-coupling parameter can be reduced to the order of \(10^{-6}\) which is much larger than the value that CMB normalization suggests for this self-coupling.

Key Words: Generalized Higgs Inflation, Cosmological Perturbations, Non-Gaussianity, Observational Constraints

PACS: 98.80.Cq , 98.80.Es

1 Introduction

Cosmological inflation is a part of the cosmic history related to a homogeneous and isotropic FRW universe that expands almost exponentially (a nearly de-Sitter universe) at very early stage of the universe evolution. The simplest model of inflation is the one in which a single scalar field with the almost flat potential runs the cosmic inflation. The theory of inflation is successful to address some problems of the standard cosmological model as well as to provide the initial density perturbations seeding the large scale structures (Guth 1981, Linde 1982, Albrecht and Steinhardt 1982, Linde 1990, Liddle and Lyth 2000, Lidsey et al. 1997, Rioto 2000, Lyth and Liddle 2009). The nearly scale invariant, adiabatic and Gaussian distribution of the perturbations modes is one of the notable predictions of the simple single field inflation (Maldacena 2003). However, by proposing some extended models of inflation and considering the non-linear perturbations, it is possible to predict some level of non-Gaussianity of the primordial perturbations (Maldacena 2003, Bartolo et al. 2004, Babich et al. 2004, Seery and Lidsey 2005, Cheung et al. 2008, Chen 2010, De Felice and Tsujikawa 2011a, De Felice and Tsujikawa 2011b, De Felice et al. 2011, Nozari and Rashidi 2012, Nozari and Rashidi 2013a, Nozari and Rashidi 2013b, Nozari and Rashidi 2014, Nozari and Rashidi 2016a, Nozari and Rashidi 2016b, Nozari and Rashidi 2017). So, it is reasonable to expect a level non-Gaussianity in future observation.

The discovery of Higgs as a fundamental particle in electro-weak symmetry breaking has significant implication in particle physics and cosmology. Regarding this fact that most of the inflation models require a scalar field (inflaton) to explain the accelerating expansion of the early universe, there is a possibility that Higgs field to be a good candidate for the inflaton (See for instance Barbon and Espinosa 2009, Calmet et al. 2017 and references therein). However, when

\(^1\)knozari@umz.ac.ir
\(^2\)s.shafizadeh@tpnu.ac.ir
\(^3\)n.rashidi@umz.ac.ir
the Higgs field is minimally coupled to the gravity, its self-coupling is too large to achieve the slow-roll inflation. Actually, to suppress the amplitude of the curvature perturbation (which should be much smaller than the Planck scale), we need to reduce the self-coupling of the Higgs field. To this end, some extensions of the Higgs inflation model have been proposed. One of these extensions is the model in which the Higgs field is non-minimally coupled to the gravity sector of the theory. In this model, a large amount of the non-minimal coupling parameter effectively suppresses the self-coupling of the Higgs field (Futamase and Maeda 1989, Salopek et al., 1989, Fakir and Urrut 1990, Kaiser 1995, Tsujikawa and Gumjudpai 2004, Bezrukov et al. 2009, Bezrukov and Shposhnikov 2008, Bezrukov and Shposhnikov 2009, Barvinsky et al., 2009, Watanabe 2011). Unfortunately, it seems that the non-minimal model violates the unitarity bound (Burgess et al. 2009, Lerner and McDonald 2010, Germani and Kehagias 2010).

It should be mentioned that the loop corrections are all small compared to the tree level amplitude and since the inflationary energy scale is always much below the scale of unitarity violation (to get constraint $\lambda \ll 1$), it is not needed to worry about the stability of inflationary model in this context (see Germani and Kehagias 2010, Calmet and Casadio 2014 for more details).

Another extension of the Higgs inflation is the new Higgs inflation model which relies on the non-minimal derivative coupling between the scalar field and Einstein tensor (Amendola 1993). By this coupling, since the normalization of the inflation field is changed, the magnitude of the Higgs self-coupling could be lower than its experimental bound.

Also in this case, due to presence of the non-renormalizable operator in the new Higgs inflationary action the time dependence unitarity bound is set. By requiring the scale of curvature that is much lower than the unitarity bound, we impose Hubble parameter scale below the Planck scale. Therefore, this postulated coupling is free of unitarity bound during inflation (Barbon and Espinosa 2009, Germani and Kehagias 2010, Atkins and Calmet 2011).

Another approach is the running kinetic inflation model in which the non-canonical kinetic term changes the normalization of the Higgs field and smoothes the general steep potential (Takahashi 2010, Dimopoulos and Thomas 2003).

The phenomenological features of the running kinetic inflation have been studied in (Nakayama and Takahashi 2008) with details. Higgs-G inflation also, is an extension of the standard Higgs inflation which incorporates the higher order derivatives of the scalar field (Kobayashi et al. 2010, Kamada et al. 2011). In the Galileon model the Lagrangian is formulated in a such way that the field equations are invariant under the Galileon symmetry $\partial_{\mu} \varphi \rightarrow \partial_{\mu} \varphi + a_{\mu}$ in the Minkowski limit of the theory (Nicolis et al. 2009, De Felice and Tsujikawa 2012). Note that, the expression $X \Box \varphi$ which is introduced as the Galileon term, emerges in the DGP model as a consequence of the combination of a branch-bending mode and a transverse graviton (Deffayet et al. 2002, Porrati 2002, Luty et al. 2003). By adding the Galileon term to the theory, the potential essentially becomes flat and the quantum fluctuations are suppressed.

In this paper, we consider another class of the generalized Galileon Higgs inflation which is a subclass of the most generalized scalar-tensor theory (Deffayet et al. 2011, Charmousis et al. 2012). Considering that the chaotic inflation is not confirmed properly by the observational data (Komatsu et al. 2010, Amsler et al. 2011, Charmousis et al. 2012), it is interesting to adopt the quartic potential and make the theory observationally viable (Germani and Kehagias 2010).

By considering an inflation model with the Galileon effect, enhanced kinetic term and non-minimal coupling between the Higgs filed and both the scalar and tensor parts of the gravity, we try to reduce the self-coupling of the Higgs sector. In doing so, we preserve also the cosmological viability of the setup. Actually the energy scale of the Higgs self-coupling constant, $\lambda$, is in the interval $0.11 < \lambda < 0.27$. From the CMB normalization, $\lambda$ is constrained to be of the order of $10^{-13}$ (Liddle and Lyth 2000)- the scale that Higgs boson can’t reach. However, by considering an inflation model with the Galileon effect, enhanced kinetic term and non-minimal coupling between the Higgs field and both the scalar and tensor parts of the gravity, we try to reduce the energy scale (self-coupling) of the Higgs sector. As we shall see, by considering this extended model, we are able to reduce the energy scale of the Higgs self-coupling constant, $\lambda$, from interval $0.11 < \lambda < 0.27$ to less than $10^{-6}$. So, in this paper, our aim is to reduce $\lambda$ by considering the Galileon-like and non-minimal effects and keeping the observational viability of the model’s parameters. In this regard, by decreasing the order of $\lambda$ and approaching the energy scale of the inflation era, the Higgs field can be considered to be an inflaton. We note that the negative values of $\lambda$ are possible in essence and at least theoretically. However, in this case inflation never happens which is out of our interest in this paper.

With these explanations, in section 2, we introduce the generalized Higgs G-inflation model and the action of the theory. In section 3 we study the background
dynamics of the model. In section 4, by adopting the ADM formalism, we expand the action up to the second and third orders of the perturbations. In this section we obtain the scalar and tensor spectral index of the primordial perturbations. We also study the non-Gaussian feature of the perturbations in both equilateral and orthogonal configurations. After that, in section 5 we perform a numerical analysis on the model’s parameter space and compare the results with Planck2015 data set. In this regard, we obtain some constraints on the model’s parameters.

2 Generalized Higgs G-Inflation Model

By detecting the Higgs boson in Large Hadron collider (LHC) experiment in Geneva (Chatrchyan et al. 2012, Aad et al. 2012), many efforts have been made to construct the inflation models where the Higgs boson acts as an inflaton. In this respect, the following Lagrangian is devoted to Higgs boson in the absence of gravity

\[ \mathcal{L}_\mathcal{H} = -D_\mu \mathcal{H}^\dagger D^\mu \mathcal{H} - \lambda (\mathcal{H}^\dagger \mathcal{H} - v^2)^2, \]

where \( D_\mu \) is the covariant derivative corresponding to the SM gauge symmetry. \( \mathcal{H} \) is the Higgs boson and \( v \sim 246 \text{ GeV} \) is its expectation value. Since the parameter \( v \) is very small compared with the Higgs field during inflation era, we can safely eliminate this parameter. Also, we concentrate on the radial part of the Higgs boson, \( \phi \sim \sqrt{2} \mathcal{H}^\dagger \mathcal{H} \), and ignore the contributions of the gauge sectors of the SM (Germani et al. 2014). In this regard, in the presence of the gravity, the Lagrangian of the Higgs model takes the following form

\[ \mathcal{L} = -m_\phi^2 R - \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{\lambda}{4} \phi^4. \]

Unfortunately, due to the large value of the Higgs self-coupling, this model in not viable (Bezrukov and Shaposhnikov 2008). Therefore, it seems reasonable to focus on the generalized G-inflation models. The action of the generalized G-inflation is written as

\[ S = \sum_{i=2}^{4} \int d^4 x \sqrt{-g} \mathcal{L}_i, \]

where \( g \) is the determinant of the metric \( g_{\mu\nu} \) and

\[ \mathcal{L}_2 = K(\phi, X), \]

\[ \mathcal{L}_3 = -G_3(\phi, X) \Box \phi, \]

\[ \mathcal{L}_4 = G_4(\phi, X) R + G_{4,X} \left( (\Box \phi)^2 - (\nabla_\mu \nabla_\nu \phi)^2 \right). \]

In the above equations, \( R = 6(\dot{H} + 2H^2) \) is the Ricci scalar, \( K \) and \( G_i \) are arbitrary functions of \( \phi \) and \( X = -(\frac{1}{2})g^{\mu\nu} \partial_\mu \phi \partial_\nu \phi \). We define \( G_i(\phi, X) = g_i(\phi) + h_i(\phi)X \). In fact, in this definition we have expanded \( G_i(\phi, X) \) as \( G_i(\phi, X) = g_i(\phi) + h_i(\phi)X + k_i(\phi)X^2 + l_i(\phi)X^3 + ... \) and just kept the terms up to the first order in \( X \) and ignored the higher order ones. We also have \( \Box \equiv g^{\mu\nu} \partial_\mu \partial_\nu \) (the standard d’Alembertian operator).

This theory has been originally found by Horndeski in a different form (Horndeski1974). This generalized model consists of the running kinetic term, the Galileon interaction, the non-minimal coupling and the non-minimal derivative coupling of the scalar field and gravity. We adopt the arbitrary functions \( K(\phi, X) \) and \( G_i(\phi, X) \) as

\[ K(\phi, X) = K(\phi)X - V, \]

\[ G_3(\phi, X) = \gamma(\phi)X, \]

\[ G_4(\phi, X) = \frac{1}{2}(m_{pl}^2 + \xi \phi^2) + \frac{1}{2\mu^2} X, \]

where \( \xi \) is a dimensionless non-minimal coupling parameter, \( \gamma(\phi) \) is a dimensionless function of the Higgs field and \( \mu \) is a mass scale. The function \( G_3 \) in equation (8) has been chosen in the way that we cover the coupling between the scalar field, kinetic term and the second-order derivatives of the scalar field (Galileon gravity). Also, in equation (9), the minimal and non-minimal coupling of the gravity with the scalar field and the coupling between the gravity and derivatives of the field have been considered. Note that, if we set \( K = 1, G_3(\phi, X) = 0 \) and \( G_4(\phi, X) = \frac{m_{pl}^2}{2} \), then the standard Higgs inflation is recovered.

3 The Background Dynamics

To derive the background equations of the model, we consider the FRW background specified by the metric

Note that, in the running kinetic inflation model, rapid growth of the kinetic term at large values of inflaton field causes the potential to be flat. In fact, in paper (Nakayama and Takahashi 2008) it has been discussed that the coefficient of the kinetic term is not necessarily unity. Actually, when the inflaton rolls over a large scale in high-scale inflation model, this coefficient is not close to 1. In this regard, to cover this issue, it is appropriate to consider the general kinetic term in the action of the model.
\[ ds^2 = -dt^2 + a^2(t)\delta_{ij}dx^i dx^j. \] By varying action (3) with respect to the metric, we find the following Friedmann equations

\[
H^2 = \frac{\rho_{\varphi}}{3m_{pl}^2}, \quad \dot{H} = -\frac{\rho_{\varphi} + p_{\varphi}}{2m_{pl}^2}, \tag{10}
\]

where the energy density and the pressure of the scalar field are defined as

\[
\rho_{\varphi} = \frac{1}{2} \dot{\varphi}^2 \left[ K + 6\gamma \frac{H \dot{\varphi}}{m_{pl}} + \gamma_{\varphi} \frac{\dot{\varphi}^2}{m_{pl}^2} + \frac{9}{2} H^2 \right] + 12 \xi \frac{\dot{\varphi} \ddot{\varphi}}{\dot{\varphi}^2} + 6 \xi \frac{H^2 \ddot{\varphi}^2}{\dot{\varphi}^2} + V, \tag{11}
\]

\[
p_{\varphi} = \frac{1}{2} \dot{\varphi}^2 \left[ K - \gamma_{\varphi} \frac{\dot{\varphi}^2}{m_{pl}^2} - 2\gamma \frac{\ddot{\varphi}}{m_{pl}} - \frac{4}{m_{pl}^2} \left( \frac{H \dot{\varphi}}{m_{pl}} \right)^2 \right] - \frac{1}{m_{pl}^2} \left( \frac{3H^2 + 2\dot{H}}{\mu^2} \right) + 4\xi + 2 \left( \frac{\dot{\varphi}^2}{\varphi^2} - 1 \right)(3H^2 + 2\dot{H}) + 4m_{pl}^2 \xi \varphi \left( \frac{\dot{\varphi}^2}{\varphi^2} + 2H \frac{\ddot{\varphi}}{\dot{\varphi}^2} \right) - V \tag{12}
\]

respectively. The equation of motion, obtained by varying the action (3) with respect to \( \varphi(t) \), is given by

\[
\frac{1}{a^3} \frac{d}{dt} (a^3 J) = P_{\varphi}, \tag{13}
\]

where

\[
J \equiv \varphi (K + \frac{3}{\mu^2} H^2) + 3\gamma H \dot{\varphi}^2 - \gamma_{\varphi} \dot{\varphi}^3, \tag{14}
\]

and

\[
P_{\varphi} = -V' - \frac{1}{2} \dot{\varphi}^2 \left( 2\gamma_{\varphi} \ddot{\varphi} + \gamma_{\varphi\varphi} \dot{\varphi}^2 - 2 \frac{\xi \dot{\varphi}}{\varphi^2} R \right). \tag{15}
\]

By substituting equations (14) and (15) into the equation (13) we get

\[
\dot{\varphi} \left( K + 6\gamma H \dot{\varphi} + \frac{3}{\mu^2} H^2 - 2\gamma_{\varphi} \dot{\varphi}^2 \right) + 3H \dot{\varphi} \left( K + 3\gamma H \dot{\varphi} + \frac{3}{\mu^2} H^2 \right) + \frac{1}{2} \dot{\varphi} \left( K_{\varphi} + 6\gamma \dot{H} - \gamma_{\varphi} \dot{\varphi}^2 \right) - 6\xi \varphi (2H^2 + \dot{H}) + V = 0. \tag{16}
\]

The slow-roll conditions in this setup are as follows

\[
\varphi^2 \ll V(\varphi), \quad |\varphi| \ll |H \dot{\varphi}|, \quad |\dot{K}| \ll |HK|, \quad |g_t(\varphi)| \ll |Hg_t(\varphi)|, \quad |\dot{h}_t| \ll |Hh_t(\varphi)|. \tag{17}
\]

By considering the slow-roll conditions, the Ricci scalar becomes

\[
R \simeq \frac{1}{m_{pl}^2 + \xi(1+6\xi)\varphi^2} \left[ 4V(\varphi) + 6\xi \varphi V' \right]. \tag{18}
\]

Also, the main background equations within the slow-roll limits take the following form

\[
H^2 \simeq \frac{V}{3(m_{pl}^2 + \xi \varphi^2)^2}, \tag{19}
\]

\[
3H \dot{\varphi} \left( K + 3H \dot{\varphi} + \frac{3}{\mu^2} H^2 \right) - \xi \dot{R} \varphi \simeq -V'. \tag{20}
\]

From equation (20) we see that in the generalized G-inflation the friction term is enhanced. By using equations (19) and (20) we can derive \( \frac{d\varphi}{dN} \equiv \frac{\dot{\varphi}}{H} \) as follows

\[
\frac{d\varphi}{dN} = \frac{\dot{\varphi}}{H} \simeq \frac{2V_{eff}'}{Y \left( K + \frac{1}{\mu^2} Y + \sqrt{(K + \frac{1}{\mu^2})Y^2 - 4\gamma V_{eff}'}^\frac{1}{4} \right)}, \tag{21}
\]

where

\[
V_{eff} = \frac{1}{m_{pl}^2 + \xi(1+6\xi)\varphi^2} \left[ -4\xi \varphi V(\varphi) + (m_{pl}^2 + \xi \varphi^2)V' \right], \tag{22}
\]

\[
Y \equiv \frac{V}{m_{pl}^2 + \xi \varphi^2}, \tag{23}
\]

and \( N = \ln a \) is the number of e-folds parameter. If the expression in the bracket of equation (22) to be small compared with the denominator, we have slow-roll inflation even with a steep potential.

The slow-roll parameters in our setup are obtained as follows

\[
\alpha = \frac{\dot{g}_t(\varphi)}{Hg_t(\varphi)} \simeq V_{eff}' \times \left( \frac{\dot{K} + (\frac{1}{\mu^2}) + \sqrt{(K + (\frac{1}{\mu^2})^2 - 4\gamma V_{eff}'})}{-4\xi \varphi V^{-1}} \right); \quad \alpha \ll 1 \tag{24}
\]

\[
\epsilon = -\frac{\dot{H}}{H^2} \simeq V^{-1} \times \left( \frac{V_{eff}'^2}{K + (\frac{1}{\mu^2}) + \sqrt{(K + (\frac{1}{\mu^2})^2 - 4\gamma V_{eff}'})^2} \right) - \frac{\alpha}{2}; \quad \epsilon \ll 1 \tag{25}
\]

\[
\eta \simeq \epsilon - \frac{1}{2\epsilon} \left( \frac{d\epsilon}{dN} \right); \quad \eta \ll 1 \tag{26}
\]
where
\[ \frac{d\epsilon}{dN} = \left( \frac{-2Y^{-1}V_{eff}'}{K + \frac{Y}{\mu^2} + \sqrt{(K + \frac{Y}{\mu^2})^2 - 4\gamma V_{eff}^2}} \right) \frac{d\epsilon}{d\varphi}, \] (27)

and
\[ \zeta = \frac{j}{HJ} \approx \epsilon - \left( \frac{V_{eff}''Y^{-1}}{K + \frac{Y}{\mu^2} + \sqrt{(K + \frac{Y}{\mu^2})^2 - 4\gamma V_{eff}^2}} \right); \quad \zeta \ll 1. \] (28)

The number of e-folds during inflation which is given by
\[ N = \int Hdt = \int \frac{H}{\dot{\varphi}}d\varphi, \] (29)

in the generalized G-inflation model takes the following form
\[ N \approx \int \left( \frac{-2\gamma Y}{K + \frac{Y}{\mu^2} + \sqrt{(K + \frac{Y}{\mu^2})^2 - 4\gamma V_{eff}^2}} \right)d\varphi. \] (30)

4 Perturbation and Non-Gaussianity

In this section, we study the perturbations in our setup. To study the tensor and scalar parts of the perturbations we should expand the action up to second order. We work in the unitary gauge ($\delta \varphi = 0$) and adopt the ADM formalism with the following metric (Baumann 2009, Mukhanov 1992)
\[ ds^2 = -N^2dt^2 + \gamma_{ij}(dx^i + N^idt)(dx^j + N^jdt), \] (31)

where $N$ and $N^i$ are the lapse and shift functions. In metric (31) we have the following definition
\[ N = 1 + 2\Phi, \quad N_i = \delta_{ij}\partial^jB, \]
\[ \gamma_{ij} = a^2(t)(1 + 2\Psi)(\delta_{ij} + h_{ij}). \] (32)

$\Phi$, $\Psi$, and $B$ are the scalar perturbations and $h_{ij}$ is the spatial shear 3-tensor. Now, we rewrite the perturbed metric up to the linear level as (Baumann 2009, Mukhanov 1992)
\[ ds^2 = -(1 + 2\Phi)dt^2 + 2a^2(t)B_ijdx^idx^j + a^2(t)(1 + 2\Psi)(\delta_{ij} + h_{ij})dx^i dx^j. \] (33)

By considering the scalar part of this metric, we expand the action (3) up to second order in the small perturbations and get
\[ S_2 = \int dtd^3x a^3 \left\{ -3 \left( \frac{m_{pl}^2 + \xi \phi^2}{a^2} \right) \dot{\Psi}^2 + \frac{1}{a^2}\left[ 2 \left( \frac{m_{pl}^2 + \xi \phi^2}{a^2} \right) \dot{\Psi} - \frac{3H(m_{pl}^2 + \xi \phi^2)}{2a^2} \right] \dot{\Phi}^2 B - \frac{2}{a^2} \left( \frac{m_{pl}^2 + \xi \phi^2}{a^2} \right) \right\} \dot{\Psi}^2 + \frac{1}{a^2}\left( \frac{m_{pl}^2 + \xi \phi^2}{a^2} \right) \right\} (\partial \Psi)^2 \]. (34)

By using the above second order action, we can find the momentum and Hamiltonian constrains as
\[ \Phi = L_1 \dot{\Psi}, \] (35)

where
\[ L_1 = \frac{2m_{pl}^2 + \xi \phi^2 - \frac{X}{\mu^2}}{H \left( 2m_{pl}^2 + \xi \phi^2 - 7 \frac{X}{\mu^2} \right)}, \] (36)

and
\[ \frac{1}{a^2}\partial^2B = 3\dot{\Psi} - \frac{1}{a^2}L_1 \partial^2 \Psi + \frac{3H(2X\phi - 7H - \frac{X}{\mu^2})}{m_{pl}^2 + \xi \phi^2} - \frac{3H^2(2X\phi - 7H - \frac{X}{\mu^2})}{m_{pl}^2 + \xi \phi^2} \dot{\Phi}, \] (37)

By substituting the equation (35) in equation (34) and integrating it by parts, the second order action reduces to the following expression
\[ S_2 = \int dtd^3x a^3 U \left[ \dot{\Psi}^2 - \frac{\xi^2}{a^2} (\partial \Psi)^2 \right]. \] (38)
Actually, there are two constraints on the sound speed smaller than (at most, equal to) the local speed of light.

\[ c \equiv 3 \left( \frac{m_{pl}^2 + \xi \phi^2 - X}{\mu^2} \right), \quad (39) \]

and

\[ c_s^2 = \frac{4 - \frac{2(\kappa + (\frac{1}{m^2})Y)}{(\kappa + (\frac{1}{m^2})Y + \sqrt{(\kappa + (\frac{1}{m^2})Y)^2 - 4\gamma V'_{eff})}}}{3 \left( \frac{m_{pl}^2 + \xi \phi^2 - X}{\mu^2} \right)}. \quad (40) \]

To avoid the ghosts and gradient instabilities it is required that

\[ U > 0, \quad c_s^2 > 0. \quad (41) \]

Actually, there are two constraints on the sound speed of the perturbations (Ellis 2007, Quiros 2017): 1- The squared sound speed of the perturbations \( (c_i^2 \) with \( i = s, T \) where \( T \) denotes the tensor part of the perturbations which we’ll study later) should be positive in order to avoid the appearance of Laplacian instabilities. That is, \( c_i^2 > 0 \). 2: From the causality requirement, the sound speed of the perturbations should be smaller than (at most, equal to) the local speed of light. That means, \( c_s^2 \leq c^2 \). Since in this paper we set \( c = 1 \), the constraint becomes \( c_s^2 \leq 1 \). These constraints are satisfied in equations (41), as we will see in the numerical analysis of non-Gaussianities. The constraints on \( c_s^2 \), lead to positive equilateral configuration of non-Gaussianities and negative orthogonal configuration of non-Gaussianities.

For convenience we define the following parameters

\[ F = 2 + \frac{1}{m_{pl}^2} \left( \xi \phi^2 + \frac{X}{\mu^2} \right), \quad (42) \]

\[ \epsilon_s = \epsilon + \frac{\phi}{m_{pl}^2 H \Phi} \left( \gamma X + \xi \phi \right). \quad (43) \]

By these definitions, we have the relation \( \epsilon_s = \frac{\mu^2}{m_{pl}^2} F \).

Causality, Laplacian and ghost free requirements impose that the right hand side of this equation to be positive. Therefore the left hand side of the relation should be positive too. On the other hand, \( \dot{H} > 0 \) means \( \epsilon < 0 \). In the case of negative \( \epsilon \), to have positive \( \epsilon_s \), the second term of equation (43) should be large enough (actually larger than \( |\epsilon| \)). Since this model is an extended one, it is likely possible to find some parameter space that gives \( \dot{H} > 0 \). However, in this paper we don’t look after this case.

The power spectrum of the curvature perturbations is given by

\[ P_\psi = \frac{H^2}{8\pi^2 U c_s^2}. \quad (44) \]

With this definition, we obtain the scalar spectral index as

\[ n_s - 1 = \frac{d \ln P_\psi}{d \ln k} \bigg|_{c_s = \epsilon_H} = -4\epsilon + \eta - \zeta + \frac{V'_{eff} \left( \kappa + (\frac{1}{m^2})Y \right)^{\prime} \left[ \kappa + (\frac{1}{m^2})Y + \sqrt{(\kappa + (\frac{1}{m^2})Y)^2 - 4\gamma V'_{eff}} \right]}{2 \left[ \kappa + (\frac{1}{m^2})Y + \sqrt{(\kappa + (\frac{1}{m^2})Y)^2 - 4\gamma V'_{eff}} \right] \left[ \kappa + (\frac{1}{m^2})Y + \sqrt{(\kappa + (\frac{1}{m^2})Y)^2 - 4\gamma V'_{eff}} \right] - 3 \left( \frac{m_{pl}^2 + \xi \phi^2 - X}{\mu^2} \right)} \]. \quad (45) \]

Now, we consider the tensor part of the metric (33) and expand the quadratic action for the tensor perturbations as follows

\[ S_T = \int dt d^3x a^3 \left[ Q \delta_{ij} - \frac{F}{a^2} (\partial h_{ij})^2 \right], \quad (46) \]

where

\[ Q \equiv (m_{pl}^2 + \xi \phi^2) - \frac{X}{\mu^2}, \quad F \equiv 2(m_{pl}^2 + \xi \phi^2) + \frac{X}{\mu^2}. \quad (47) \]

The sound speed square is given by

\[ c_T^2 = \frac{F}{Q} = \frac{m_{pl}^2 + \xi \phi^2 + \frac{X}{\mu^2}}{m_{pl}^2 + \xi \phi^2 - \frac{X}{\mu^2}}. \quad (48) \]

Note that, satisfying conditions \( Q > 0 \) and \( c_T^2 > 0 \) lead to the ghost and Laplacian free perturbations. We note that constraint from observation of gravitational waves by LIGO/VIRGO opens a research area at this point.
The power spectrum of primordial tensor perturbations is given by
\[ \mathcal{P}_T = \frac{H^2}{2\pi^2 Qc_T^2} \simeq \left( \frac{1}{2m_{pl}^2 + 2\xi\varphi^2 + \frac{\lambda}{\mu^2}} \right) \frac{2H^2}{\pi^2}, \] (49)
leading to the following tensor spectral index
\[ n_T = \frac{d\ln\mathcal{P}_T}{d\ln k} \simeq \left( \frac{V_{eff}'}{V} \right)^2 \times \left( \frac{m_{pl}^2 \xi \varphi^2}{\mathcal{K} + \left( \frac{1}{\mu^2} \right) Y + \sqrt{\left( \mathcal{K} + \left( \frac{1}{\mu^2} \right) Y \right)^2 - 4\gamma V_{eff}'}} \right). \] (50)

The tensor-to-scalar ratio in this setup is given by
\[ r = \frac{\mathcal{P}_T}{\mathcal{P}_\psi} \simeq \frac{8}{3\sqrt{3}} \left( \frac{4 - \frac{2(\mathcal{K} + \left( \frac{1}{\mu^2} \right) Y)}{\mathcal{K} + \left( \frac{1}{\mu^2} \right) Y + \sqrt{\left( \mathcal{K} + \left( \frac{1}{\mu^2} \right) Y \right)^2 - 4\gamma V_{eff}'}}}{2 - \frac{2(\mathcal{K} + \left( \frac{1}{\mu^2} \right) Y)}{\mathcal{K} + \left( \frac{1}{\mu^2} \right) Y + \sqrt{\left( \mathcal{K} + \left( \frac{1}{\mu^2} \right) Y \right)^2 - 4\gamma V_{eff}'}}} \right)^{\frac{1}{2}} n_T. \] (51)

As we see, the consistency relation is modified in the presence of Galileon effect. If we set \( G_3 = 0, G_4 = \frac{m_{pl}^2}{2} \) and \( \mathcal{K}(\varphi) = 1 \) in equation (51), the model recovers the standard consistency relation \( r = -8n_T \).

By regarding this fact that for a Gaussian distribution, any odd point correlation functions vanishes, to seek for the non-Gaussian feature we should study three point correlation function (Ohashi et al. 2013). To this end, we expand action up to the third order in the small perturbations. We eliminate parameter \( B \) by using equation (35) and introduce \( (\chi) \) as
\[ B = -L_1 \Psi + \frac{a^2 \chi}{(2m_{pl}^2 + \xi \varphi^2 - \frac{\lambda}{\mu^2})}, \]
\[ \chi = \left( \frac{2m_{pl}^2 + \xi \varphi^2 - \frac{\lambda}{\mu^2}}{a^2} \right) B. \]

In this regard, we obtain the cubic action as
\[ S_3 = \int dt \mathcal{L}_3, \] (52)
where
\[ \mathcal{L}_3 = \int d^3x \left\{ a^3 \left( \frac{F_{\epsilon}^2}{c_s^2} \left( -3 \left( \frac{1}{c_s^2} - 1 \right) + \frac{1}{c_s^2} \left( \epsilon - \frac{\dot{\epsilon}}{H_{\epsilon s}} \right) \right) \right. \right.
\[ + \frac{\varphi}{m_{pl}^2 HF} \left( \xi \varphi - 3\gamma X \right) - \frac{6}{m_{pl}^2 \mu^2 F} \right) \left( m_{pl}^2 \Psi \dot{\Psi}^2 \right) + a \left( F_{\epsilon s} \left( \frac{1}{c_s^2} - 1 \right) + F_{\epsilon s} \left( \epsilon_s + \frac{\dot{\epsilon}}{H_{\epsilon s}} - 2 \frac{\dot{\epsilon}}{H_{\epsilon s}} \right) \right) \left( \frac{1}{2m_{pl}^2 + \xi \varphi^2 + \frac{\lambda}{\mu^2}} \right) \left( \frac{2X}{\mu^2} \right) \left( m_{pl}^2 \Psi (\partial \Psi)^2 \right) + a^3 \left( \frac{m_{pl}^2 F_{\epsilon s}}{H_{\epsilon s}^3} \left( \frac{1}{c_s^2} - 1 - 2 \frac{\Lambda}{\Sigma} \right) + \frac{1}{c_s^2} \left( \frac{\gamma X \varphi}{m_{pl}^2 HF} \right) \right.
\[ \left. + \frac{2}{m_{pl}^2 \mu^2 F} - \frac{\xi \varphi \dot{\varphi}}{m_{pl}^2 HF} \right) - 3 \frac{\gamma X \varphi}{m_{pl}^2 HF} + \frac{\xi \varphi \dot{\varphi}}{m_{pl}^2 HF} \right) \left. \frac{2}{m_{pl}^2 \mu^2 F} - \frac{6 \left( \frac{\gamma X \varphi}{m_{pl}^2 HF} \right)}{c_s^2} \right) \left( m_{pl}^2 \Psi^3 - 2a^2 \frac{\epsilon_s}{c_s^2} \right)
\[ \left. \ddot{\Psi} \left( \partial_i \Psi \right) \left( \partial_i \chi \right) + a^3 \left( \frac{1}{4m_{pl}^2} \left( \epsilon_s - 4 \frac{\gamma X \varphi}{m_{pl}^2 HF} \right) \right) \left. \left[ \ddot{\Psi} \left( \partial_i \Psi \right)^2 - \partial_i \partial_j \left( \partial_i \chi \right) \right] + \left. a \left( \frac{2 \gamma X \varphi}{HF^2} \right) \left( \ddot{\Psi} \partial_i \Psi \partial_j \chi - \partial_i \partial_j \left( \partial_i \chi \right) \right) \right) \right\}. \] (53)

In this equation, parameters \( \Sigma \) and \( \Lambda \) are defined as
\[ \Sigma \equiv \left( \frac{m_{pl}^2 F - 2 \frac{\chi}{\mu^2}}{m_{pl}^4} \right) \left[ \left( m_{pl}^2 HF - \gamma X \varphi - 4 \frac{H X}{\mu^2} \right) \right.
\[ + m_{pl}^2 \xi \varphi \dot{\varphi} \left. \right)^2 + \left( m_{pl}^2 F - 2 \frac{X}{\mu^2} \right) \left( -3 m_{pl}^2 H^2 F + X \right)
\[ + 12 \gamma H X \dot{\varphi} - 4 \gamma \varphi X^2 + \frac{21 H X}{\mu^2} - 6 m_{pl}^2 (H \varphi \dot{\varphi}) \right), \] (54)
and
\[ \Lambda \equiv \left( 2 + \frac{1}{m_{pl}^2} \left( \xi \varphi^2 + \frac{X}{2m_{pl}^2} \right) \right) \left( \gamma H X \varphi - \frac{4}{3} \gamma \varphi X^2 \right). \] (55)

To obtain the three point correlators, we should calculate the vacuum expectation value of the curvature perturbations during inflation as follows (see for instance (Maldacena 2003, Cheung et al. 2008, Seery and Lidsey 2005))
\[ \langle \Psi(k_1) \Psi(k_2) \Psi(k_3) \rangle = -i \int_{\tau_i}^{\tau_f} d\sigma a(0) \langle \Psi(0, k_1) \Psi(0, k_2) \Psi(0, k_3), \mathcal{H}_{int}(\tau) \rangle |0 \rangle, \] (56)
where interacting Hamiltonian is $\mathcal{H}_{int} = -\mathcal{L}_3$. We can assume that the dimensionless coefficient of each contribution in the third order action can be treated as a constant because of the slow varying of those coefficients during the inflation epoch. In this respect, by solving the integral (56) we get

$$\langle \Psi_{k_1} \Psi_{k_2} \Psi_{k_3} \rangle = (2\pi)^3 \delta(k_1 + k_2 + k_3) B_\Psi(k_1, k_2, k_3),$$

with

$$B_\Psi = \frac{(2\pi)^4}{4 \prod_{i=1}^3 k_i^3} (P_\Psi)^2 A_\Psi(k_1, k_2, k_3).$$

The resulting bispectrum is achieved by considering that the additional shape functions can be defined by using other shape functions (mentioned in (Renaux-Petel 2012, De Felice and Tsujikawa 2013)) in the Horn-deski’s theories, which can be written as

$$A_\Psi(k_1, k_2, k_3) = \left\{ \frac{3}{2} \left( \frac{1}{c_s^2} - 1 - \frac{2A}{\Sigma} + \frac{6}{c_s^2} \left( \frac{\gamma X \dot{\varphi}}{m_{pl}^2 HF} \right) \right) - \frac{6}{c_s^2} \left( \frac{\gamma X \dot{\varphi}}{m_{pl}^2 HF c_s} \right) \left( \prod_{i=1}^3 k_i^2 \right) \right\} \left( \frac{2}{K} \sum_{i>j} k_i^2 k_j^2 \right) - \frac{4}{K} \sum_{i>j} k_i^2 k_j^2 + \frac{2}{K^2} \sum_{i<j} k_i^2 k_j^3 \right\} (59)

The following parameter gives the amplitude of the non-Gaussianity

$$f_{NL} = \frac{10}{3} \sum_{i=1}^3 k_i^3.$$

Following Refs. (Renaux-Petel 2012, De Felice and Tsujikawa 2013) we introduce the following shapes

$$S_{ortho} = \frac{18}{13} \left\{ \left( \frac{2}{K} \sum_{i>j} k_i^2 k_j^2 \right) - \frac{1}{K} \sum_{i\neq j} k_i^2 k_j^3 \right\} - \frac{216}{13} \left( \frac{\prod_{i=1}^3 k_i^2}{K^3} \right).$$

and

$$S_{ortho} = \frac{12}{14 - 13\beta} \left[ (3 - \frac{9}{2} \beta) \left( \frac{2}{K} \sum_{i>j} k_i^2 k_j^2 - \frac{1}{K} \sum_{i\neq j} k_i^2 k_j^3 \right) + \frac{3}{2} \beta - 1 \left( \sum_i k_i^2 - \frac{4}{K^2} \sum_{i\neq j} k_i^2 k_j^3 \right) \right] + 18\beta \left( \frac{\prod_{i=1}^3 k_i^2}{K^3} \right),$$

which are orthogonal. Now, we rewrite equation (59) in terms of $S_{equil}$ and $S_{ortho}$ (Renaux-Petel 2012) as

$$A_\Psi = a_1 S_{equil} + a_2 S_{ortho},$$

where

$$a_1 = \frac{13}{12} \left[ \frac{1}{24} \left( \frac{1}{c_s^2} - 1 \right) (2 + 3\beta) \right] + \frac{A}{12\Sigma} (2 - 3\beta) \right] - \frac{1}{6\epsilon_s} \left( \frac{\gamma X \dot{\varphi}}{m_{pl}^2 HF} \right) (2 - 3\beta) + \frac{1}{3\epsilon_s^2} \left( \frac{\gamma X \dot{\varphi}}{m_{pl}^2 HF} \right) \right].$$

and

$$a_2 = \frac{14 - 13\beta}{12} \left[ \frac{1}{8} \left( \frac{1}{c_s^2} - 1 \right) - \frac{A}{4\Sigma} + \frac{1}{2\epsilon_s} \left( \frac{\gamma X \dot{\varphi}}{m_{pl}^2 HF} \right) \right].$$

We can obtain the amplitudes of the non-Gaussianity in the equilateral and orthogonal configurations from equations (60)-(62) as follows

$$f_{NL}^{equil} = \left[ \frac{1}{24} \left( \frac{1}{c_s^2} - 1 \right) (2 + 3\beta) \right] + \frac{A}{12\Sigma} (2 - 3\beta) \right] - \frac{1}{6\epsilon_s} \left( \frac{\gamma X \dot{\varphi}}{m_{pl}^2 HF} \right) (2 - 3\beta) + \frac{1}{3\epsilon_s^2} \left( \frac{\gamma X \dot{\varphi}}{m_{pl}^2 HF} \right).$$

and

$$f_{NL}^{ortho} = \left[ \frac{1}{8} \left( \frac{1}{c_s^2} - 1 \right) - \frac{A}{4\Sigma} + \frac{1}{2\epsilon_s} \left( \frac{\gamma X \dot{\varphi}}{m_{pl}^2 HF} \right) \right].$$

Considering that at $k_1 = k_2 = k_3$ limit, both the equilateral and orthogonal configurations have a maximal signal, we obtain the non-linear parameters in this
limit. The results are as
\[
\begin{align*}
{f_{NL}}_{\text{equil}} &= \frac{325}{18} \left( \frac{1}{24} \left( 1 - \frac{1}{c_s^2} \right) (2 + 3\beta) + \frac{\Lambda}{12\Sigma} (2 - 3\beta) \right),
\end{align*}
\]
\[
- \frac{1}{6\epsilon_s} \left( \frac{\gamma X \dot{\phi}}{m_{pl}^2 \hat{H} \tilde{F}} \right) (2 - 3\beta) + \frac{1}{3\epsilon_s c_s^2} \left( \frac{\gamma X \dot{\phi}}{m_{pl}^2 \hat{H} \tilde{F}} \right),
\]
and
\[
{f_{NL}}_{\text{ortho}} = \frac{10}{9} \left( \frac{65}{4} \beta + \frac{7}{6} \right) \left[ \frac{1}{8} \left( 1 - \frac{1}{c_s^2} \right) - \frac{\Lambda}{4\Sigma} + \frac{1}{2\epsilon_s} \left( \frac{\gamma X \dot{\phi}}{m_{pl}^2 \hat{H} \tilde{F}} \right) \right],
\]
(68)
(69)
After calculation of perturbations and possible non-Gaussianity of these perturbations we compare our results with observations in the next section.

5 Confrontation with Observational Data

In this section we perform a numerical analysis on the parameter space of our generalized G-inflation model and compare the results with Planck2015 observational data. To this end, we adopt a potential as \( V = \frac{1}{4} \varphi^4 \) and we set \( K = 1 \) and \( \gamma(\varphi) = \varphi/M^4 \). To perform the numerical analysis we assume \( M \approx 10^{-5} m_{pl} \), \( \mu \approx 3 \times 10^{-8} m_{pl} \) and \( N = 60 \). Now, by solving the integral of equation (30), we obtain the value of the Higgs field at the horizon crossing of the physical scales. After that, by using this obtained value we can find the scalar spectral index, tensor-to-scalar ratio and the amplitudes of the equilateral and the orthogonal configurations of the non-Gaussianity in terms of \( N \), \( \lambda \) and \( \xi \). Then, we analyze the model parameter space numerically. The results are shown in figures.

Figure 1 shows the ranges of the self-coupling parameter of the Higgs field, \( \lambda \), and the non-minimal coupling parameter, \( \xi \), that lead to the observationally viable values of the scalar spectral index and tensor-to-scalar ratio. In plotting the figures we have focused on \( \lambda < 10^{-6} \) and \( \xi < 2 \times 10^2 \). Figure shows that, as \( \xi \) increases the smaller values of \( \lambda \) are observationally viable. In figure 2 we have plotted the tensor-to-scalar ratio versus the scalar spectral index in the background of Planck2015 TT, TE, EE+lowP data. To plot this figure, we have adopted three sample values of the non-minimal coupling parameter as \( \xi = 50, \xi = 80 \) and \( \xi = 100 \). Our numerical analysis shows that this generalized G-inflation model is consistent with Planck2015 data if \( 10^{-7} \leq \lambda \leq 2 \times 10^{-6} \) for \( \xi = 50 \), \( 10^{-7} \leq \lambda \leq 5 \times 10^{-6} \) for \( \xi = 80 \) and \( 10^{-7} \leq \lambda \leq 7 \times 10^{-6} \) for \( \xi = 100 \). Note that the presence of the Galileon-like interaction and the NMC effect in this model cause a reduction of the tensor-to-scalar ratio in comparison to the standard situation. We have also studied the amplitudes of the non-Gaussianity in both the equilateral and orthogonal configurations numerically. The results are shown in figures 3 and 4. We have analyzed \( {f_{NL}}_{\text{equil}} \) and \( {f_{NL}}_{\text{ortho}} \) in the ranges of the parameters used in studying \( n_s \) and \( r \). Figures 3 and 4 show that in the ranges \( \lambda < 10^{-6} \) and \( \xi < 2 \times 10^2 \), both equilateral and orthogonal non-Gaussianities are consistent with Planck2015 TT, EE, TTE and EET data. As these figures show, in this generalized G-inflation model, it is possible to have large non-Gaussianity in some subspaces of the model parameter space. From our analysis we can say that if we consider a generalized G-inflation model, depending on the values of \( \xi \), it is possible to have \( \lambda < 10^{-6} \) (specially, \( \lambda \sim 10^{-13} \) which is well in the range of CMB result (Liddle and Lyth 2000)). This means that, if we adopt smaller values of \( \xi \), it is possible to reduce the self-coupling of the Higgs sector in order to reach the energy scale of inflation in this setup. This is an important results since it provides a possible mechanism for reduction of the Higgs self-coupling as an inflaton.

6 Summary

In this paper we have studied the cosmological inflation in a generalized G-inflation model. We have studied the effects of the Galileon interaction and the non-minimal coupling on the energy scale of the Higgs inflation. In this regard we have adopted the non-minimal coupling function as \( \xi \varphi^2 \) and other functions as \( V = \frac{1}{4} \varphi^4, K = 1 \) and \( \gamma(\varphi) = \frac{\varphi}{M^4} \). We have obtained the background dynamics and then we have treated the perturbations in this generalized setup in details. By expanding the action up to the second order, we have obtained the scalar and tensor spectral indices and tensor-to-scalar ratio in this generalized G-inflation model. In this respect, we have shown that the presence of the Galileon effect modifies the consistency relation. By calculating the cubic action and the three point correlation function, we have studied the non-Gaussian feature of perturbations in this setup. We have also obtained the nonlinear parameters in both equilateral and orthogonal configurations of the non-Gaussianity at \( k_1 = k_2 = k_3 \) limit. Finally, we have performed a numerical analysis on the model’s parameter space to obtain some constraints on the parameters. We have studied \( n_s \), \( r \), \( f_{NL}^{\text{equil}} \) and \( f_{NL}^{\text{ortho}} \) numerically. Our numerical analysis shows that if we consider the non-minimal coupling
Fig. 1  Ranges of $\lambda$ and $\xi$ leading to the observationally viable values of the scalar spectral index (left panel) and the tensor-to-scalar ratio (right panel) for a generalized Higgs Galileon model. We note that consistency with observations in this generalized model requires enhancement of $\lambda$.

Fig. 2  Tensor-to-scalar ratio versus the scalar spectral index for a generalized Higgs model, in the background of Planck2015 TT, TE, EE+lowP data. The figure is plotted with N=60.
Fig. 3 Ranges of $\lambda$ and $\xi$ leading to observationally viable values of the amplitudes of the equilateral (left panel) and orthogonal (right panel) configurations of the non-Gaussianity for a generalized Higgs inflation model. In both panels all the adopted ranges are consistent with Planck 2015 observational data.

Fig. 4 Amplitude of the orthogonal configuration of the non-Gaussianity versus the amplitude of the equilateral configuration for a generalized Higgs inflation in the background of Planck 2015 TTT, EEE, TTE and EET data. The figure is plotted with N=60. Note that the diagrams for all three values of the non-minimal coupling are too close to be distinguished from each other in this figure.
and Galileon-like interactions, it is possible to control the values of the self-coupling parameter $\lambda$. Actually, in this extended model, depending on the values of the non-minimal coupling, we were able to reduce the values of $\lambda$ from interval $0.11 < \lambda < 0.27$ to $\lambda < 10^{-6}$. In fact, if we adopt smaller values of $\xi$, it is possible to reduce the energy scale (self-coupling) of the Higgs sector in order to reach the energy scale of inflation ($\lambda \sim 10^{-13}$) in this setup. Therefore, by reducing the order of $\lambda$ and approaching the energy scale of the inflation era, the Higgs field can be considered to be an inflaton.

Acknowledgements
The work of K. Nozari has been supported financially by Research Institute for Astronomy and Astrophysics of Maragha (RIAAM) under research project number 1/4717-71.
References

Aad G. et al.: Phys. Lett. B 716, 1 (2012).
Albrecht A. and Steinhardt P.: Phys. Rev. D 48, 1220 (1993).
Amendola L.: Phys. Lett. B 301, 175 (1993).
Amelino C. et al.: Phys. Lett. B 667, 1 (2008).
Atkins C. and Calmet X.: Phys. Lett. B 697, 37 (2011).
Babich D., Creminelli P. and Zaldarriaga, M.: J. Cosmol. Astropart. Phys., 0408, 009 (2004).
Barbon J. L. F. and Espinosa J. R.: Phys. Rev. D 79, 083502 (2009).
Barolo N., Komatsu E., Matarrese S. and Riotto A.: Phys. Rept.: 402, 103, (2004).
Barvinsky A. O. et al.: J. Cosmol. Astropart. Phys. 0912, 003 (2009).
Baumann, D.: [arXiv:hep-th/0907.5424].
Bezrukov F. and Shaposhnikov M.: J. Cosmol. Astropart. Phys. 0906, 029 (2009).
Bezrukov F. and Shaposhnikov, M.: Phys. Lett. B 659, 703 (2008).
Bezrukov F. and Shaposhnikov M.: J. High Energy Phys. 0907, 089 (2009).
Burgess C. P., Lee H. M. and Trott, M.: J. High Energy Phys. 0909, 103 (2009).
Calmet X. and Casadio R., Phys. Lett. B 734, 17 (2014).
Calmet X. et al., [arXiv:1701.02140]
Charmousis C. et al.: Phys. Rev. Lett. 108, 051101 (2012).
Chatrchyan S. et al.: Phys. Lett. B 716, 30 (2012).
Chen X.: Adv. Astron. 2010, 638979 (2010).
Cheung C. et al.: J. High Energy Phys., 0803, 014, (2008).
Cheung C. et al.: J. High Energy Phys., 0803, 014, (2008).
Ellis E. et al.: Gen. Rel. Grav., 39, 1651, (2007).
De Felice A. and Tsujikawa S.: Phys. Rev. D 84, 083504, (2011).
Dimopoulos S. and Thomas S. D.: Phys. Lett. B 573, 13 (2003).
De Felice A. et al.: J. Cosmol. Astropart. Phys. 1108, 021 (2011).
De Felice A. and Tsujikawa S.: J. Cosmol. Astropart. Phys. 1202, 007 (2012).
De Felice A. and Tsujikawa S.: J. Cosmol. Astropart. Phys. 03, 030 (2013).
Deffayet C. et al.: Phys. Rev. D 65, 044026 (2002).
Deffayet C. et al.: Phys. Rev. D 84, 064039 (2011).
Fakir R. and Unruh W. G.: Phys. Rev. D 41, 1783 (1990).
Futamase T. and Maeda K. i.: Phys. Rev. D 39, 399 (1989).
Germani C. and Kehagias A.: Phys. Rev. Lett. 105, 011302 (2010).
Germani C., Watanabe Y. and Wintergerst N.: J. Cosmol. Astropart. Phys., 12, 009, (2014).
Guth A.: Phys. Rev. D 23, 347 (1981).
Horndeski G. W.: Int. J. Theor. Phys. 10, 363, (1974).
Kaiser D. I.: Phys. Rev. D 52, 4295 (1995).
Kamada K., et. al: Phys. Rev. D 83, 083515 (2011).
Kobayashi T. et. al: Phys. Rev. Lett. 105, 231302 (2010).
Komatsu, E. et al.: [arXiv:astro-ph/1001.4538].
Lerner R. N. and McDonald J.: J. Cosmol. Astropart. Phys. 1004, 015 (2010).
Liddle A. and Lyth, D.: Cosmological Inflation and Large-Scale Structure, Cambridge University Press, (2000).
Liddle J. E. et al.: Rev. Mod. Phys. 69, 373, (1997).
Linde A.: Phys. Lett. B 108 , 389 (1982).
Linde, A.: Particle Physics and Inflationary Cosmology Harwood Academic Publishers, Chur, Switzerland, (1990), [arXiv:hep-th/0503203].
Luty M. A. et al.: J. High Energy Phys. 0309, 029 (2003).
Lyth D. H. and Lididdle, A. R.: The Primordial Density Perturbation, Cambridge University Press, (2009).
Maldacena J. M.: J. High Energy Phys., 0305, 013, (2003).
Mukhanov V. F. et al.: Phys. Rep. 215, 203 (1992).
Nakayama K. and Takahashi, F.: [arXiv:hep-ph/1008.2956].
Nicolas A. et al.: Phys. Rev. D 79, 064036 (2009).
Nozari K. and Rashidi N.: Phys. Rev. D 86, 043505 (2012).
Nozari K. and Rashidi N.: Phys. Rev. D 88, 023519 (2013).
Nozari K. and Rashidi N.: Phys. Rev. D 88, 084040 (2013).
Nozari K. and Rashidi N.: Astrophys. Space Sci. 350, 339 (2014).
Nozari K. and Rashidi N.: Phys. Rev. D 93, 124022 (2016).
Nozari K. and Rashidi N.: Advances in High Energy Physics, 2016, 1252689 (2016).
Nozari K. and Rashidi, N.: [arXiv:1705.02617].
Ohashi J., Soda J. and Tsujikawa S.: J. Cosmol. Astropart. Phys. 1312, 009 (2013).
Porrati M.: Phys. Lett. B 534, 209 (2002).
Quiros I., Gonzalez T., Nucamendi U., Garca-Salcedo R., Horta-Rangel F. A. and Saavedra J ., [arXiv:1707.03885] [gr-qc].
Renaux-Petel S.: J. Cosmol. Astropart. Phys. 1202, 020 (2012).
Riotto A. [arXiv:hep-ph/0210162].
Salopek D. S. et al.: Phys. Rev. D 40, 1753 (1989).
Seery D. and Lidsey J. E.: J. Cosmol. Astropart. Phys. 0506, 003 (2005).
Takahashi F.: Phys. Lett. B 693, 140 (2010).
Tsujikawa S. and Gumjudpai B.: Phys. Rev. D 69, 123523 (2004).
Watanabe Y.: Phys. Rev. D 83, 043511 (2011).

This manuscript was prepared with the AAS L\LaTeX macros v5.2.