Revisiting slow-roll dynamics and the tensor tilt in general single-field inflation

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We explore the possibility of a blue-tilted gravitational wave spectrum from potential-driven slow-roll inflation in the Horndeski theory. In Kamada et al. (2012), it was claimed that a blue gravitational wave spectrum cannot be obtained from stable potential-driven slow-roll inflation within the Horndeski framework. However, it has been demonstrated that the spectrum of primordial gravitational waves can be blue in inflation with the Gauss-Bonnet term, where the potential term is dominant and slow-roll conditions as well as the stability conditions are satisfied. To fill in this gap, we clarify where the discrepancy is coming from. We extend the formulation of Kamada et al. (2012) and show that a blue gravitational wave spectrum can certainly be generated from stable slow-roll inflation if some of the conditions previously imposed on the form of the free functions in the Lagrangian are relaxed.

I. INTRODUCTION

Inflation [1–3] is the most promising candidate of the scenario for the early Universe, explaining naturally the large-scale homogeneity of the observed Universe and the origin of primordial density perturbations that lead to the CMB fluctuations and the large-scale structure. Inflation also predicts the existence of a stochastic background CMB fluctuations and the large-scale structure. Inflation of primordial gravitational waves (tensor perturbations), which has yet to be observed. In the most standard inflationary scenario, a quasi-de Sitter expansion is driven by the (rather conservative) case of potential-driven slow-roll inflation within the Horndeski framework. However, it has been demonstrated that the spectrum of primordial gravitational waves can be blue in inflation with the Gauss-Bonnet term, where the potential term is dominant and slow-roll conditions as well as the stability conditions are satisfied. To fill in this gap, we clarify where the discrepancy is coming from. We extend the formulation of Kamada et al. (2012) and show that a blue gravitational wave spectrum can certainly be generated from stable slow-roll inflation if some of the conditions previously imposed on the form of the free functions in the Lagrangian are relaxed.

While the argument in [9] seems to be general to a large extent, a counterexample is known to exist in the literature: inflation with the nonminimal coupling to the Gauss-Bonnet term [18–29]. Though in this case the energy density of the slowly rolling inflaton is dominated by its potential, one can have a positive tensor tilt [18, 19, 22]. Since the nonminimal coupling between a scalar field and the Gauss-Bonnet term is just a specific example of the Horndeski Lagrangian [4], there must be something overlooked in the analysis of [9].

The purpose of this paper is to fill in the gap between the above apparently contradicting statements. In fact, the formulation of [9] is not general enough to accommodate Gauss-Bonnet inflation. Moreover, an unnecessarily strong assumption was made in [9]. In this paper, we improve these points and enlarge a possible model space of slow-roll inflation within the Horndeski theory, showing that blue gravitational waves can indeed be generated from (stable) slow-roll inflation.

This paper is organized as follows. In the next section, we review the previous study [9] and suggest a possible improvement as implied by the example of Gauss-Bonnet inflation. In Sec. III, we extend the slow-roll dynamics to cover the inflationary model space which has not been explored in [9]. We then discuss the possibility of a blue-tilted gravitational wave spectrum from potential-driven slow-roll inflation in Sec. IV. Finally, we draw our conclusion in Sec. V.

II. SLOW-ROLL INFLATION FROM HORNDESKI

A. A quick recap of Kamada et al. [9]

Let us review briefly the argument of Ref. [9], where generic slow-roll inflationary dynamics is investigated. The analysis of Ref. [9] is based on the Horndeski theory, i.e., the most general scalar-tensor theory with second-order field equations, whose Lagrangian is given by

\[ \mathcal{L} = G_2(\phi, X) - G_3(\phi, X)\Box \phi + G_4(\phi, X)R \]

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Here, $\phi$ is the scalar field, $X := -g^\mu \nabla_\mu \phi \nabla_\nu \phi / 2$, $R$ is the Ricci scalar, and $G_{\mu\nu}$ is the Einstein tensor. The background cosmological equations and the quadratic action governing cosmological perturbations in the Horndeski ground cosmological equations and the quadratic action reduce to \[9\] under these conditions reduce to \[9\] and their spectral index, is the Hubble parameter and a dot denotes differentiation with respect to the cosmic time. With some manipulation, the background cosmological equations under these conditions reduce to \[9\]

$$
6g_4H^2 \simeq V, \quad -4g_4H + 2g_4H \simeq \dot{\phi} \left( u + 3vH\dot{\phi} \right), \quad 3H\dot{\phi} \simeq \frac{1}{2v} \left( -u + \sqrt{u^2 - 4Uv} \right),
$$

where we defined

$$
u(\phi) := h_2 + \frac{1}{2v}, \quad v(\phi) := h_3 + \frac{h_5 V}{6g_4},$$

and a prime denotes differentiation with respect to $\phi$. It is obvious that Eq. (4) is essentially the Friedmann equation. Equations (5) and (6) correspond respectively to the familiar equations $-2M^2_{Pl}H = \dot{\phi}^2$ and $3H\dot{\phi} \simeq -V'$ in the canonical slow-roll inflation model.

In Ref. \[9\] it is assumed that

$$u(\phi) > 0,$$

probably because $u$ determines the sign of the kinetic term of $\phi$ in the simple case with $h_4 = 0$ and hence is expected to be correlated with some of the stability conditions. This is another assumption which we revisit carefully in this paper.

Using Eqs. (5) and (6) we obtain

$$2\epsilon + \delta_M = \frac{\dot{\phi}^2}{4g_4H^2} \left( u + \sqrt{u^2 - 4Uv} \right).$$

It follows from the assumption (8) that

$$2\epsilon + \delta_M > 0.$$

This inequality will lead to the important conclusion on the tensor tilt.

The quadratic action for tensor perturbations in generic potential-driven inflation is given by \[9\]

$$S^{(2)}_h = \frac{1}{4} \int dt dz dz^3 x a^3 g_4 \left[ \frac{\dot{h}_t^2}{\dot{\phi}^2} - a^2 (\partial_i h_t)^2 \right].$$

A nonstandard feature appears only in the time-dependent effective Planck mass $g_4$. The stability condition is equivalent to the aforementioned assumption $g_4 > 0$. Following the usual quantization procedure one obtains the tensor power spectrum,

$$P_T = \frac{H^2}{\pi^2 g_4},$$

and its spectral index,

$$n_t = -2\epsilon - \delta_M.$$

Thus, from Eq. (10) we see that the tensor power spectrum would never be blue in generic potential-driven inflation,
Using Eq. (6) we obtain
\[ F = \frac{\dot{\phi}^2}{6H^2} \left( u + 2\sqrt{u^2 - 4U'v} \right), \] (18)
\[ \epsilon_s^2 = \frac{1}{3} \left( 2 + \frac{u}{\sqrt{u^2 - 4U'v}} \right). \] (19)

Thus, under the assumption (8), the stability conditions \( F > 0 \) and \( c_s^2 > 0 \) are indeed satisfied. However, \( u > 0 \) seems to be only a sufficient condition for the stability.

It is straightforward to calculate the power spectrum and the spectral index [4]:
\[ P_\zeta = \frac{1}{8\pi^2 c_s F}, \] (20)
\[ n_s - 1 = -2\epsilon - \frac{\dot{c_s}}{Hc_s} - \frac{\dot{F}}{HF}. \] (21)

B. Possible improvement of Kamada et al. [9]: The case of Gauss-Bonnet inflation

It was pointed out that the tensor spectral index can be positive in (stable) slow-roll inflation with the Gauss-Bonnet term [18, 19, 28]. More explicitly, in Refs. [18, 19, 28] the following nonminimal coupling to the Gauss-Bonnet term is considered:
\[ f(\phi) \left( R^2 - 4R_{\mu\nu}R^{\mu\nu} + R_{\mu\nu\rho\sigma}R^{\mu\nu\rho\sigma} \right). \] (22)

It is well-known that this term yields second-order field equations, and hence resides within the Horndeski theory. Nevertheless, the result of Refs. [18, 19, 28] seems to be inconsistent with the conclusion of Ref. [9]. This indicates that the work of [9], which was probably built upon somewhat stronger assumptions than necessary, can potentially be improved to accommodate potential-driven inflation models with a blue tensor spectrum such as Gauss-Bonnet inflation.

Let us start with looking at how the nonminimal coupling to the Gauss-Bonnet term (22) is incorporated into the Horndeski theory. As shown in Ref. [4], the way of reproducing the term (22) from the Horndeski functions is nontrivial:
\[ G_2 \supset 8f'''X(3 - \ln X), \quad G_3 \supset 4f'''X(7 - 3\ln X), \]
\[ G_4 \supset 4f'''X(2 - \ln X), \quad G_5 \supset -4f'\ln X. \] (23)

This clearly shows that the Taylor-expanded form (2) fails to capture the structure of the Gauss-Bonnet term. The first three terms could be slow-roll suppressed because they are proportional to second or higher derivatives of weakly \( \phi \)-dependent function \( f \), but \( G_5 \supset -4f'\ln X \) cannot be ignored even in the slow-roll regime. This observation hints at how we can proceed to extend the framework of [9].

Moreover, we have seen that the previous assumption (8) is likely to be too strong. It is therefore desirable to revisit this point and clarify the necessary condition for the stability. This will also enlarge the possible model space of slow-roll inflation explored by Ref. [9].

III. SLOW-ROLL DYNAMICS

To capture the essential part of the Gauss-Bonnet term in the slow-roll regime, let us now assume that the Horndeski functions take the form of
\[ G_a = g_a(\phi) + \lambda_a(\phi) \ln X + h_a(\phi)X + \cdots, \] (24)
where the ellipsis stands for slow-roll suppressed terms. As in the previous analysis, we eliminate \( g_3 \) and \( g_5 \) by performing integration by parts, and assume that \( g_4 > 0 \). One may also consider the terms of the form \( \xi_a(\phi)X \ln X \), which could be as large as, or even larger than, \( h_a X \). In the present analysis, however, we will assume that \( \lambda_a \) is already of first order in the slow-roll approximation, \( \lambda_a = \mathcal{O}(h_a X) \), and accordingly \( \xi_a X \ln X \) is of second order. The newly introduced terms \( \lambda_a \ln X \) would be dangerous in the \( X \to 0 \) limit. However, as we will see below, at least some of them yield only regular terms at the level of field equations.

We assume the same slow-roll conditions as given in Eq. (3). For the new functions we impose the analogous slow-roll conditions,
\[ \frac{\dot{\lambda}_a}{H\lambda_a} \ll 1. \] (25)

Under these assumptions, the time-time component of the gravitational field equations for a cosmological background reduces to
\[ 6g_4H^2 \simeq V + \lambda_2(2 - \ln X) - 6\lambda_4H^2\ln X \]
\[ + 6\lambda_3H\dot{\phi} + 6\lambda_5H^3\dot{\phi}. \] (26)

To avoid the singular terms in the \( X \to 0 \) limit, we require that
\[ \lambda_2 = 0, \quad \lambda_4 = 0. \] (27)

As seen from Eq. (26), \( \lambda_3 \) and \( \lambda_5 \) do not lead to any singular terms in the field equations, and hence are acceptable. In addition to the slow-roll conditions, we impose the following potential-dominance conditions on these two functions:
\[ \delta_3 := \frac{\lambda_3 H \dot{\phi}}{g_4 H} \lesssim \mathcal{O}(\epsilon), \quad \delta_5 := \frac{\lambda_5 H \dot{\phi}}{g_4} \lesssim \mathcal{O}(\epsilon), \] (28)

namely, the last two terms in the right-hand side of Eq. (26) are actually of the same order of the other slow-roll suppressed terms. We thus have the same equation as Eq. (4),
\[ 6g_4H^2 \simeq V, \] (29)
even in the presence of the \( \ln X \) terms. This equation rules out the possibility of \( g_4 < 0 \) (for \( V > 0 \)).

The space-space components of the gravitational field equations and the equation of motion for the scalar field in the slow-roll regime reduce to
\[ -4g_4 \dot{H} + 2g_4 H \simeq \dot{\phi}^2 \mathcal{I}, \] (30)
\[ 3H \dot{\mathcal{I}} \simeq -U'(\phi), \] (31)
where
\[ I := u(\phi) + 3v(\phi)H\dot{\phi} + \varpi(\phi) \left( \frac{V}{3H^2} \right), \]  
(32)
\[ \varpi(\phi) := \varpi_3(\phi) + \varpi_5(\phi), \]  
(33)
\[ \varpi_3(\phi) := \frac{3V}{g_4} \lambda_3, \quad \varpi_5(\phi) := \frac{1}{6} \left( \frac{V}{g_4} \right)^2 \lambda_5, \]  
(34)
and \( u, v, \) and \( U' \) were defined earlier in Eq. (7). Now it is easy to solve Eq. (31) for \( 3H\dot{\phi} \) to get
\[ 3H\dot{\phi} \simeq \frac{1}{2v} \left[ -u \pm \sqrt{u^2 - 4v(U' + \varpi)} \right]. \]  
(35)
At this stage we have two branches (if \( v \neq 0 \)), but it will turn out in the end that the “−” branch exhibits instabilities of scalar perturbations. Therefore, here we only consider the “+” branch. Using Eq. (35), one can rewrite \( I \) in terms of the functions of \( \phi \) as
\[ I = \frac{U'}{2(U' + \varpi)} \left[ u + \sqrt{u^2 - 4v(U' + \varpi)} \right]. \]  
(36)
If \( v = 0 \), we do not need to care about the branches and we instead have
\[ 3H\dot{\phi} \simeq -\frac{U'}{u}, \]  
(37)
\[ I = \frac{U' u}{U' + \varpi}. \]  
(38)
Equation (30) can also be expressed as
\[ 2c + \delta_M = \frac{\dot{\phi}^2 I}{2g_4 H^2}. \]  
(39)
This is the generalization of Eq. (9) and will be used later in the next section.

Let us take a look at the role of \( I \) in the slow-roll dynamics, focusing on the simple case with \( g_4 = M_P^2/2 \). Using the background equations (29), (30), and (31), the potential slow-roll parameter, \( c_V := M_P^2 (V'/V)^2/2 \), can be written as
\[ \epsilon = \epsilon_V I. \]  
(40)
This implies that, even if the potential is too steep to support usual inflation (say, \( \epsilon_V = \cal{O}(1) \)), inflation can still occur provided that \( I \gg 1 \).

To highlight the impact of the newly introduced term, let us further focus on the case with \( u = 1 \) and \( v = 0 \). In this case, the Lagrangian is given by \( \mathcal{L} = (M_P^2/2)R + X + V + (\ln X \text{ corrections}) \). We then have
\[ I = \frac{V'}{V' + \varpi} \Rightarrow \epsilon = \frac{M_P^2}{2} \frac{V'(V' + \varpi)}{V^2}, \]  
(41)
\[ 3H\dot{\phi} \simeq -(V' + \varpi). \]  
(42)
Thus, \( \varpi \) effectively shifts the potential slope. If the potential is nearly flat and \( \varpi \lesssim \cal{O}(V') \), the \( \ln X \) terms have only a minor effect on the dynamics. In the opposite limit, \( \varpi \gg \cal{O}(V') \), inflation is spoiled. The most interesting case is that \( V' \) could be large but is canceled by \( \varpi \); \( V' \simeq -\varpi \). In this case, inflation occurs with the help of the \( \ln X \) terms.

It should be emphasized that so far we have made no assumption about the sign of \( u \). Only the constraint coming from the background dynamics is that the expression in the square root must be nonnegative:
\[ u^2 \geq 4v(U' + \varpi). \]  
(43)

IV. TENSOR TILT AND STABILITY
Let us move to the main question of this paper: Is a blue tensor spectrum compatible with the potential-driven slow-roll dynamics and the stability of cosmological perturbations?

We substitute the assumed form of the Horndeski functions \( G_a \) [Eq. (24)] to the general formulas of the quadratic action for cosmological perturbations derived in Ref. [4]. We then make the slow-roll and potential-dominance approximations. Even if one takes into account the \( \ln X \) terms in \( G_a \), it is found that under these approximations the quadratic action for tensor perturbations remains the same as Eq. (11).

The tensor spectral index is given by
\[ n_t = -2\epsilon - \delta_M = -\frac{\dot{\phi}^2 I}{2g_4 H^2}, \]  
(44)
where we used Eq. (39). Thus, the sign of \( I \) plays the key role in determining the sign of \( n_t \).

The quadratic action for the curvature perturbation takes the form of Eq. (15), but now with
\[ F = \frac{\dot{\phi}^2}{2H^2} \left( u + 4vH\dot{\phi} + \frac{4\varpi_3}{9H^2} \right), \]  
(45)
\[ c_s^2 = \frac{u + 4vH\dot{\phi} + 4\varpi_3/9H^2}{u + 6vH\dot{\phi}}. \]  
(46)
For the stability of the scalar sector it is necessary that both denominator and numerator of \( c_s^2 \) are positive. Using Eq. (35), we have
\[ u + 6vH\dot{\phi} = \begin{cases} \pm \sqrt{u^2 - 4v(U' + \varpi)} & (v \neq 0) \\ u & (v = 0) \end{cases}. \]  
(47)
Thus, as long as we take the “+” branch, one of the stability conditions is automatically satisfied. However, in the case of \( v = 0 \), the stability condition requires \( u > 0 \).

The numerator of \( c_s^2 \) reads
\[ u + 4vH\dot{\phi} + 4\varpi_3/9H^2 \]  
\[ = \begin{cases} \frac{(1 - 2A)}{3} u + \frac{2(1 - A)}{3} \sqrt{u^2 - 4v(U' + \varpi)} & (v \neq 0) \\ \left(1 - \frac{4A}{3}\right) u & (v = 0) \end{cases}. \]  
(48)
where $A := \varpi_3/(U' + \varpi)$. Since the several independent functions participate in the stability conditions, it is not so illuminating to analyze the most general case. Instead let us consider the following two special cases: (i) $\varpi_3 = \varpi_5 = 0$, and (ii) $v = 0$.

(i) $\varpi_3 = \varpi_5 = 0$

Since the ln $X$ terms vanish in $G_\alpha$, this case corresponds to the reanalysis of Ref. [9]. The tensor tilt and the stability are determined by the two combinations of the functions, $u$ and $vU'$. The stability condition reads

$$u + 2\sqrt{u^2 - 4vU'} > 0,$$

while the tensor tilt depends on the sign of

$$I = \frac{1}{2} \left( u + \sqrt{u^2 - 4vU'} \right).$$

We plot in Fig. 1 the stable region in the $(u, vU')$ plane. It is found that the scalar perturbations are stable and $n_t > 0$ if

$$u < 0, \quad 0 < vU' < \frac{3}{16} u^2.$$  \hspace{1cm} (51)

To realize $n_t > 0$, it is not necessary to extend the formulation of [9]. The region (51) was just overlooked in the previous analysis.

(ii) $v = 0$

In this case, one can easily see how the ln $X$ terms help to realize $n_t > 0$. The stability conditions reads

$$u > 0 \quad \text{and} \quad \frac{\varpi_3}{U' + \varpi_3 + \varpi_5} < \frac{3}{4}. \hspace{1cm} (52)$$

From Eq. (38) and $u > 0$ we see that the tensor tilt depends on the sign of $U'/\left(U' + \varpi\right)$. We have a blue tensor spectrum if

$$\frac{U'}{U' + \varpi_3 + \varpi_5} < 0. \hspace{1cm} (53)$$

Equations (52) and (53) are satisfied simultaneously if

$$\frac{\varpi_3}{U'} - 3\frac{\varpi_5}{U'} > 3 \quad \text{and} \quad \varpi_3 + \varpi_5 < -1. \hspace{1cm} (54)$$

Figure 2 summarizes the stability and the tensor tilt in the $(\varpi_3/U', \varpi_5/U')$ plane. Note that from the definitions and Eqs. (31) and (39) we see that $\varpi_3/U' = -3\delta_M/(2\epsilon + \delta_M)$ and $\varpi_5/U' = -2\delta_M/(2\epsilon + \delta_M)$, and hence we typically have $|\varpi_3/U'|, |\varpi_5/U'| \lesssim O(1)$.

Let us finally check that the result of Gauss-Bonnet inflation studied in [18, 19, 28] can be reproduced from our general framework. The Lagrangian considered in [28] is given by\footnote{In terms of the notation of Koh et al. [28], $f(\phi) \rightarrow -\xi(\phi)/2$. Note also that $\epsilon_{\text{Koh}} = \epsilon$, $\eta_{\text{Koh}} = -2\epsilon + \dot{\epsilon}/H\epsilon$, $\delta_{\text{Koh}} = \delta$, $\delta_{2\text{Koh}} = \epsilon + \delta/\dot{\delta}$.}

$$\mathcal{L} = \frac{M_{Pl}^2}{2} R + X - V(\phi) + f(\phi) \left( R^2 - 4R_{\mu\nu}R^{\mu\nu} + R_{\mu\nu\rho\sigma}R^{\mu\nu\rho\sigma} \right). \hspace{1cm} (55)$$

![FIG. 1. Tensor tilt and stability in the reanalysis of [9]. The $u < 0$ region was overlooked in the previous analysis. Note in passing that the white region does not satisfy the condition (43).](image1)

![FIG. 2. Tensor tilt and stability in $v = 0$ models.](image2)
Ignoring the higher-order terms in the slow-roll approximation, this corresponds to
\begin{align}
G_2 &= -V + X, \quad G_3 = 0, \quad G_4 = \frac{M_{Pl}^2}{2}, \\
G_5 &= -4f' \ln X,
\end{align}
and hence
\begin{align}
\dot{u} = 1, \quad v = 0, \quad U' = V', \quad \varpi = \varpi_5 = -\frac{8}{3M_{Pl}^2}.
\end{align}

It follows from Eq. (54) that stable inflation with \( n_t > 0 \) is realized if
\begin{align}
\frac{\varpi_5}{U'} < -1 \Leftrightarrow \frac{f'}{M_{Pl}^2 V'} > \frac{3}{8}.
\end{align}
This reproduces the result of [28].

The expression for the scalar spectral index can also be reproduced. Using Eq. (30), we obtain \( \dot{\varphi}^2 = M_{Pl}^2 H^2 (2\epsilon - \delta_5) \), and hence \( F = M_{Pl}^2 (\epsilon - \delta_5/2) \). This, together with Eq. (21), leads to
\begin{align}
n_s - 1 = -2 \epsilon - \frac{(2\epsilon - \delta_5)}{(2\epsilon - \delta_5)H},
\end{align}
which agrees with the result of [28]. In our notation, this can also be written in a simpler way as \( n_s - 1 = -4\epsilon + 2\eta \). (See Ref. [27] for a higher-order calculation of the power spectra in Gauss-Bonnet inflation.)

We have thus successfully extended the work of [9] to include Gauss-Bonnet inflation as a specific case. It should be emphasized that our framework not only contains Gauss-Bonnet inflation but also covers a model space of potential-driven slow-roll inflation which has not been explored before.

V. CONCLUSIONS

In this paper, we have explored the possibility of generating primordial gravitational waves with a positive tilt, \( n_t > 0 \), from potential-driven slow-roll inflation. In a canonical inflationary setup this is obviously impossible. While a blue-tipped gravitational wave spectrum can certainly be obtained in kinetically driven G-inflation [10] and other more or less radical models [12–17], it has been claimed in [9] that one always has \( n_t < 0 \) in generic potential-driven slow-roll inflation constructed within the Horndeski framework [4–6]. It was pointed out, however, that the tensor tilt can indeed be positive in potential-driven slow-roll inflation with the Gauss-Bonnet term [18, 19, 28]. Since the nonminimal coupling to the Gauss-Bonnet term is a (nontrivial) specific case of the Horndeski Lagrangian, this result implies that the validity of the statement of [9] is questionable. In this work, we have therefore extended the formulation of [9] in two ways. First, the Taylor-expanded form of the functions in the Horndeski Lagrangian assumed in [9] fails to reproduce inflation with the Gauss-Bonnet term and so we have included new terms that are still allowed within the Horndeski framework and help to recover the Gauss-Bonnet term as a specific case. Second, we have reconsidered the validity of the inequality assumed in [9] and found that it is in fact too strong. We have shown that a blue gravitational wave spectrum can be obtained even in potential-dominated slow-roll inflation if one relaxes at least either of these two assumptions. We have thus enlarged a possible model space of slow-roll inflation with observationally interesting predictions.

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