I. INTRODUCTION

Inflation is accepted as a testable paradigm for the physics of the early Universe. This scenario addresses some important problems of the standard big bang cosmology satisfactorily. The most important problem in standard cosmology is the origin of the cosmological structures. In this regard, the most favorite inflationary models are those models in which the needed primordial density perturbations, seeding the structure formation in the Universe, can be generated naturally \[1, 2, 3, 4, 5, 6, 7, 8\]. In the simplest realization of the inflation, the Universe is dominated by a nearly flat potential energy of a single scalar field (the inflaton field). Quantum fluctuations of the scalar field during the inflationary epoch are considered as the primordial source of the cosmological perturbations. The dominant mode of the primordial density perturbations in the simple single field model (with canonical kinetic term) is almost adiabatic and scale invariant. Also, the distribution of the perturbations for this simple case is predicted to be Gaussian \[9\]. Nevertheless, with the advancement of technology, the observational cosmologists have shown that the primordial density perturbations are somehow scale dependent \[10, 11\]. Despite this fact that currently there is no direct signal for non-Gaussian feature of the perturbations, Planck collaboration has obtained some constraints on the primordial non-Gaussianity \[12\]. In this regard, by proposing extended inflationary models, some authors have tried to realize theoretically a level of non-Gaussianity in the dominant mode of the primordial density perturbation \[13, 14, 15, 16, 17, 18, 19, 21, 22, 23, 24, 25, 26\].

Carrying a large amount of information about the cosmological dynamics by the non-Gaussian perturbations, makes the models predicting non-Gaussianity more interesting.

Motivated by \(\alpha\)-attractor models, in this paper we consider a Gauss-Bonnet inflation with an E-model-type of potential. We consider the Gauss-Bonnet coupling function to be the same as the E-model potential. In the small \(\alpha\) limit we obtain an attractor at \(r = 0\) as expected, and in the large \(\alpha\) limit we recover the Gauss-Bonnet model with potential and coupling function of the form \(\phi^2\). We study perturbations and non-Gaussianity in this setup and we find some constraints on the model’s parameters in comparison with PLANCK data sets. We study also the reheating epoch after inflation in this setup. For this purpose, we seek the number of e-folds and temperature during reheating epoch. These quantities depend on the model’s parameter and the effective equation of state of the dominating energy density in the reheating era. We find some observational constraints on these parameters.

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Key Words: Inflation, Cosmological Perturbations, Non-Gaussianity, Reheating, \(\alpha\)-Attractor, Observational Constraints
made the assumption that the inflaton field, as a scalar field responsible for cosmological inflation, belongs to the dilaton fields family.

Recently, the idea of “cosmological attractors” in the inflation models has attracted much attention. Among the models incorporating the idea of cosmological attractors, we can refer to conformal attractors [43, 44] and α-attractors models [47, 48, 49, 50]. For more details on the issue of α-attractor see also Refs. [51, 52, 53, 54, 55, 56, 57, 58]. The conformal attractor models have the characteristic property that in large N limit, the spectral index of the primordial curvature perturbations and the tensor-to-scalar ratio always are given as $n_s = 1 - \frac{1}{3} \alpha$ and $r = \frac{12 \alpha}{N^2}$ respectively. For a single field α-attractor models, in the small α limit we have $n_s = 1 - \frac{1}{3} \alpha$ and $r = \frac{12 \alpha}{N^2}$. In this paper, we are interested in α-attractor models. There are two types of this models called E-model and T-model. The E-model is specified by the potential as

$$V = V_0 \left[ 1 - \exp \left( -\sqrt{\frac{2\kappa^2}{3\alpha}} \phi \right) \right]^{2n}, \quad (1)$$

and the T-model is specified by

$$V = V_0 \tanh^{2n} \left( \frac{\kappa \phi}{\sqrt{\alpha \Omega}} \right), \quad (2)$$

where $V_0$, $n$ and $\alpha$ are some free parameters. In the large $\alpha$ limit, the predictions of these models converge towards the predictions of the simple single field inflation model with the potential as $V \sim \phi^{2n}$. In this paper we are interested in E-models of α-attractors and study primordial perturbations and non-Gaussianity in this class of models.

An important issue in cosmological inflation is the reheating process after the end of inflation. As long as the slow-roll conditions ($\eta, \epsilon \ll 1$) (requiring the potential to be sufficiently flat) are satisfied, the Universe inflates. As soon as the slow-roll conditions break down, the inflaton rolls to the minimum of its potential and starts to oscillate. In a simple canonic reheating scenario, this oscillation of the inflaton, by processes including the physics of particle creation and non-equilibrium phenomena, causes the inflaton decay to the plasma of the relativistic particles leading to the radiation-dominated Universe [59, 60, 61]. Nevertheless, some other complicated scenarios of reheating have been proposed which include also the non-perturbative processes. Among them we can refer to instant preheating [62], the parametric resonance decay [63, 64, 65] and tachyonic instability [66, 67, 68, 69, 70, 71]. The reheating era, is characterized by parameters $N_{rh}$ (number of e-folds during reheating) and $\omega_{eff}$ (the effective equation of state parameter during reheating). In this regard, studying $N_{rh}$, $\omega_{eff}$ and also $T_{rh}$ (the reheating temperature) in the inflation models seems to be useful to constrain these models [72, 73, 74, 75, 76]. See also Ref. [77] for a recent and elegant review on reheating.

In this paper, we consider an inflation model in which the scalar field is nonminimally coupled to the Gauss-Bonnet (GB) term in 4-dimension. The potential of the scalar field is considered to be E-model α-attractor.

$$G(\phi) = G_0 \left[ 1 - \exp \left( -\sqrt{\frac{2\kappa^2}{3\alpha}} \phi \right) \right]^{2n}. \quad (3)$$

In this regard, when $\alpha \to \infty$, the GB coupling function approaches to $G(\phi) \sim G_0 \phi^{2n}$ and when $\alpha \to 0$, the GB term has no effect on the dynamics. In section II, by obtaining the main equations, we study the background dynamics of this GB-α-attractor model. In section III we study the linear perturbations (including the scalar and tensor perturbations) in our setup and we obtain expressions for the scalar spectral index and tensor-to-scalar ratio in terms of $\alpha$ and $G_0$. Then we perform the numerical analysis on the model’s parameter space. In section VI, the non-linear perturbations and non-gaussian feature of perturbations distribution are considered in details. We obtain the amplitude of the non-Gaussianity in terms of the sound speed (which is a function of the model’s parameters) and in the equilateral limit of the momenta.

Then we study the evolution of $f_{NL}$ versus $c_s^2$ numerically. After that, in section V, we study the reheating epoch in this GB-α-attractor model. In this section we calculate expressions for the number of e-folds and temperature during the reheating phase after inflation. In this regard we obtain more constraints on the model’s parameters. Finally, in section VI, the summary of our analysis is presented.

II. THE SETUP

The action for an inflation model including a nonminimally coupled Gauss-Bonnet term in 4-dimensions is written as

$$S = \int d^4 x \sqrt{-g} \left[ \frac{1}{2\kappa^2} R - \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) - G(\phi) \mathcal{L}_{GB} \right], \quad (4)$$

where, $R$ is the Ricci scalar, $\phi$ is an inflaton filed and $V(\phi)$ is the E-model potential defined in equation (1). Also, $G(\phi)$ is the Gauss-Bonnet coupling term given by equation (3). $\mathcal{L}_{GB}$ is the lagrangian of the Gauss-Bonnet term. For a spatially flat FRW geometry, Friedmann equation of this model is given as follows

$$H^2 = \frac{\kappa^2}{3} \left( \frac{1}{2} \dot{\phi}^2 + V(\phi) + 2\kappa H^3 \dot{G} \right), \quad (5)$$

where a dot marks a time derivative of the corresponding parameter. The following equation of motion

$$-\ddot{\phi} - 3H \dot{\phi} - 4\kappa^2 H^4 G' - 24H^2 \dot{H} G' - V' = 0 \quad (6)$$
is obtained by varying the action \( I \) with respect to the inflaton field and a prime marks differentiation with respect to \( \phi \).

The definition of the slow-roll parameters as \( \epsilon = -\frac{\dot{\phi}^2}{2H^2} \) and \( \eta = -\frac{\dot{\phi}^2}{H^2} \), gives the following expressions for \( \epsilon \) and \( \eta \) in terms of the model’s parameters

\[
\epsilon = -\frac{\mathcal{F}}{2H - 24H^2\kappa^2\dot{\phi}^2} H^2, \\
\eta = \frac{\dot{\mathcal{F}}}{H\mathcal{F}} + \frac{2\dot{H} - 48H\kappa^2\dot{\phi}^2 H - 24H^2\kappa^2\dot{\phi}^2}{H(2H - 24H^2\kappa^2\dot{\phi}^2)},
\]

where the parameter \( \mathcal{F} \) is defined as follows

\[ \mathcal{F} \equiv 8H^3\kappa^2\dot{\phi}^2 + \frac{\kappa^2}{3}V'\dot{\phi} + \frac{\kappa^2}{3}\dot{\phi}\ddot{\phi}. \]

Since the Hubble parameter evolves during the inflationary era so slowly, the conditions \( \epsilon \ll 1 \) and \( \eta \ll 1 \) are satisfied in this period. When one of these parameters reaches unity, the inflation phase terminates automatically.

Although in a simple single field model the slow-roll limits are defined as \( \ddot{\phi} \ll |3H\dot{\phi}| \) and \( \dot{\phi}^2 \ll V(\phi) \), in the presence of GB term there are more conditions. To have the slow-roll regime in this model, the conditions \( 8\kappa^2 H|\dot{\phi}| \ll 1 \) and \( |\dot{\phi}^2| \ll |\dot{\phi}|H | \) should also be satisfied (see [21, 39]). By regarding these conditions and considering the slow evolution of the Hubble parameter \( \dot{H} \ll H^2 \), we get

\[ H^2 \simeq \frac{\kappa^2}{3}V, \]

and

\[ -3H\ddot{\phi} - 24H^4\dot{G}' - V' \simeq 0. \]

In this model, from the definition of the e-folds number as

\[ N = \int_{t_{hc}}^{t_f} H dt, \]

where \( t_{hc} \) and \( t_f \) are the time of horizon crossing and end of inflation respectively, and equation \( [1] \), we get

\[ N \simeq \int_{\phi_{hc}}^{\phi_f} -\frac{3H^2}{V' + 24H^2\dot{G}'}d\phi. \]

In this paper, we are going to investigate the cosmological viability of the GB model with E-model type potential. In this regard, one way is to study the perturbations and non-Gaussian features of this model. In the following section, we focus on the linear perturbations to obtain the tensor and scalar spectral indices and ratio of their amplitudes.

### III. LINEAR PERTURBATIONS

By using the Arnowitt-Deser-Misner (ADM) formalism which is given by the following metric \[ 78 \]

\[ ds^2 = -N^2dt^2 + h_{ij}(dx^i + N^i dt)(dx^j + N^j dt), \]

where \( N^i \) is the shift vector and \( N \) is the lapse function, one can expand the action \( I \) up to the second and third orders in the small fluctuations of the space-time background metric (which contribute in the perturbations). By expanding the shift and laps functions in terms of the 3-scalars \( \mathcal{R} \) and \( \mathcal{Y} \) and a vector \( v^i \) (with condition \( v^i = 0 \) as \( N^i \equiv \mathcal{Y} = \delta_{ij}\partial_j\mathcal{Y} + v^i \) and \( N = 1 + \mathcal{R} \)), the general perturbed form of the ADM metric \[ 14 \] can be obtained \[ 79, 80 \]. By defining the parameters \( \Theta_{ij} \) as a spatial symmetric and traceless shear 3-tensor and \( D \) as the spatial curvature perturbation, we can write the coefficient \( h_{ij} \) in the second term of equation \( [14] \) as \( h_{ij} = a^2[(1 - 2D)\delta_{ij} + 2\Theta_{ij}] \). Therefore, we can write the perturbed form of the metric \( [13] \) as

\[ ds^2 = -(1 + 2\mathcal{R})dt^2 + 2a(t)\mathcal{Y}_i dt dx^i \\
+ a^2(t)(1 - 2D)\delta_{ij} + 2\Theta_{ij}) dx^i dx^j. \]

Working within the uniform-field gauge (specified by \( \delta\phi = 0 \)) is convenient to study the scalar perturbation of the theory. In this regard, by adopting \( \Theta_{ij} = 0 \) and considering the scalar part of the metric components, we get the following expression for the perturbed metric \[ 79, 80, 51 \]

\[ ds^2 = -(1 + 2\mathcal{R})dt^2 + 2a(t)\mathcal{Y}_i dt dx^i \\
+ a^2(t)(1 - 2D)\delta_{ij} dx^i dx^j. \]

The second order (quadratic) action is obtained by replacing the perturbed metric \[ 10 \] in action \[ 4 \] and expanding it up to the second order in small perturbations. The result is as follows

\[
S_2 = \int dt d^3x a^3 \left\{ -3(\kappa^{-2} - 8\mathcal{G})\mathcal{D}^2 - \frac{2(\kappa^{-2} - 8\mathcal{G})}{a^2} \mathcal{R}\mathcal{D}^2 \mathcal{D} + \frac{1}{a^2} \left[ 2(\kappa^{-2} - 8\mathcal{G})\mathcal{D} - (\kappa^{-2} H - 24H^2\mathcal{G})\mathcal{R} \right] \partial^2 \mathcal{Y} + 3\left( 2\kappa^{-2} H - 24H^2\mathcal{G} \right) \mathcal{R} \mathcal{D} - \left( 3\kappa^{-2} H^2 - \frac{\dot{\mathcal{G}}}{2} - 48H^3\mathcal{G} \right) \mathcal{R}^2 \right. \\
\left. + \frac{\kappa^{-2} - 8H\mathcal{G}}{a^2}(\partial\mathcal{D})^2 \right\}. \]

The following expressions are the equations of motion of \( \mathcal{R} \) and \( \mathcal{Y} \), which are obtained by variation of the quadratic action \[ 17 \] with respect to the corresponding parameters

\[ \mathcal{R} = 2\left( \frac{\kappa^{-2} - 8H\mathcal{G}}{2\kappa^{-2} H - 24H^2\mathcal{G}} \right), \]
If we use the equation (18) to replace \( \mathcal{R} \) in action (17) and taking some integration by part, we get (to see more details on calculation of the second and third order action, see Refs. [15, 16, 21, 23, 82])

\[
S_2 = \int dt \, d^3 x \, \mathcal{W}_s \left[ \dot{D}^2 - \frac{c_s^2}{a^2} (\partial D)^2 \right],
\]

where the parameters \( \mathcal{W}_s \) and \( c_s^2 \) (dubbed sound velocity), are functions of the model’s parameters as follows

\[
\mathcal{W}_s \equiv -4 \left( \kappa^{-2} - 8H \dot{\varphi} \right)^2 \left( 9\kappa^{-2}H^2 - \frac{3}{2} \dot{\varphi}^2 - 144H^3 \dot{\varphi} \right) \quad \frac{3}{2} \left( 2\kappa^{-2}H - 24H^2 \dot{\varphi} \right)^2 + 3 \left( \kappa^{-2} - 8H \dot{\varphi} \right), \]

and

\[
c_s^2 = 3 \left[ 2 \left( 2\kappa^{-2}H - 24H^2 \dot{\varphi} \right) \left( \kappa^{-2} - 8H \dot{\varphi} \right) H - \left( 2\kappa^{-2}H - 24H^2 \dot{\varphi} \right)^2 \left( \kappa^{-2} - 8H \dot{\varphi} \right)^{-1} \left( \kappa^{-2} - 8\dot{\varphi} \right) \right]
+ 4 \left( 2\kappa^{-2}H - 24H^2 \dot{\varphi} \right) \frac{d \left( \kappa^{-2} - 8H \dot{\varphi} \right)}{dt} - 2 \left( \kappa^{-2} - 8H \dot{\varphi} \right) \frac{d(2\kappa^{-2}H - 24H^2 \dot{\varphi})}{dt} \left[ \left( 9 \left( 2\kappa^{-2}H - 24H^2 \dot{\varphi} \right)^2 - 4 \left( \kappa^{-2} - 8H \dot{\varphi} \right) \left( 9\kappa^{-2}H^2 - \frac{3}{2} \dot{\varphi}^2 - 144H^3 \dot{\varphi} \right) \right) \right]^{-1}.
\]

When we survey the perturbations of a model, studying the scalar spectral index gives some useful information about the primordial perturbations. To obtain this parameter, we should find the power spectrum in the model, by calculating the vacuum expectation value of \( D \)

\[
(0|D(0,k_1)D(0,k_2)|0) = (2\pi)^3 \delta^3(k_1 + k_2) \frac{2\pi^2}{k^3} A_s.
\]

By \( \tau = 0 \), we mean the time of the end of inflationary era. In equation (23), the power spectrum \( A_s \) is defined as

\[
A_s = \frac{H^2}{8\pi^2 \mathcal{W}_s c_s^2}.
\]

Now, we are in the position to obtain the scalar spectral index from the power spectrum as follows

\[
n_s - 1 = \left. \frac{d \ln A_s}{d \ln k} \right|_{c_s k = aH} = -2\epsilon - \frac{1}{H} \frac{d \ln(\epsilon - 4\kappa^2 H\dot{\phi})}{dt} - \frac{1}{H} \frac{d \ln c_s}{dt}.
\]

By calculating the value of the scalar spectral index, we can seek the scale dependence of the primordial perturbations.

As the scalar part of the perturbations, studying the tensor part also gives some other information about the initial perturbations. In this regard, we consider the tensor part of the perturbed metric [15]. We set the shear 3-tensor \( \Theta_{ij} \) to be as

\[
\Theta_{ij} = \Theta_{+} \delta_{ij} + \Theta_{\times} \delta_{ij}^\times,
\]

which is written in terms of the two polarization tensors. These polarization tensors satisfy the reality and normalization conditions [15] [16]. The following expression is the quadratic action for the tensor mode of the perturbations (namely, gravitational waves)

\[
S_T = \int dt \, d^3 x \, a^3 W_T \left[ \dot{\Theta}_+^2 - \frac{c_T^2}{a^2} (\partial \Theta_+)^2 + \dot{\Theta}_\times^2 - \frac{c_T^2}{a^2} (\partial \Theta_\times)^2 \right].
\]

The parameters \( W_T \) and \( c_T^2 \) in the second order action [27] are expressed as follows

\[
W_T \equiv \frac{1}{4\kappa^2} \left( 1 - 8\kappa^2 H \dot{\varphi} \right),
\]

\[
c_T^2 \equiv 1 + 8\kappa^2 H \dot{\varphi}.
\]

To obtain the tensor spectral index \( (n_T) \), we should firstly obtain the amplitude of the tensor perturbations which is given in our setup as

\[
A_T = \frac{H^2}{2\pi^2 W_T c_T^2},
\]
which gives
\[ n_T = \frac{d \ln A_T}{d \ln k} = -2\epsilon. \tag{31} \]
Finally, tensor-to-scalar ratio is obtained as
\[ r = \frac{A_T}{A_s} \simeq 16c_s \left( \epsilon - 4\kappa^2 H G' \right). \tag{32} \]
This equation shows that the presence of the Gauss-Bonnet coupling modifies the standard consistency relation which in a canonical single field model is given by \( r = 16 \) (or \( r = -8n_T \)). We note that in some other inflation models such as the non-minimal gravitational coupling models and multiple fields inflation the consistency relation is modified too. \cite{12, 54, 85}.

For a numerical analysis on the perturbations parameters \( n_s \) and \( r \), we adopt the potential and the GB coupling function as defined in equations \( \ref{eq:1} \) and \( \ref{eq:3} \) (the E-model). In this regard, we obtain the following expressions
\[ n_s = 1 - \frac{3\alpha V^2}{8n^2\kappa^4U^2} - \left( \frac{\alpha \dot{V}}{8n^2\kappa^4U^2} \right) \sqrt{\frac{\kappa^2}{\alpha}} + \frac{\sqrt{6\alpha} V \dot{\phi}}{8n^2\kappa^4U^2} \sqrt{\frac{\kappa^2}{\alpha}} \]
\[ - \frac{64\zeta^3}{3n^2} V_0 \nu^{4n-2} G_n U_0 \nu^{2n} \phi^2 + \frac{16\alpha}{3n^2} V_0 \nu^{4n-1} G_n U_0 \nu \phi \int d\phi \]
\[ \left( \frac{3\alpha V^2}{16n^2\kappa^4U^2} - \frac{8}{3n^2} V_0 \nu^{4n-1} G_n U_0 \nu \phi \right)^{-1} \frac{\kappa^2}{\alpha} \right] V_0^{-1} \nu^{-2n} \]
\[ - \frac{d \ln c_s^2}{d t} V_0^{-1} \nu^{-2n}, \tag{33} \]
where we have defined the parameters \( V \) and \( U \) as follows
\[ V = 1 - U, \tag{35} \]
\[ U = e^{-\frac{2\phi}{\kappa^2}}. \tag{36} \]

We take \( n = 1 \) and study the evolution of the tensor-to-scalar ratio versus the scalar spectral index numerically. The results are shown in figure 1. We have plotted \( r \) versus \( n_s \) for five values of the nonminimal coupling parameter \( G_0 \). When \( \alpha \to 0 \) all trajectories converge to the point shown by the red star, corresponding to \( r \to 0 \) and \( n_s \to 0.967 \). Also, when \( \alpha \to \infty \), the trajectories tend to the point corresponding to the ones with \( V \sim \frac{2}{\kappa^2} \) and \( G \sim G_0 \frac{\phi^2}{\kappa^2} \). This GB-\( \alpha \)-attractor model is consistent with observational data in some ranges of \( G_0 \) and \( \alpha \). In this manner we constrain \( G_0 \) for some values of \( \alpha \) in comparison with Planck2015 \cite{10} observational data. The results are summarized in table I.

IV. NONLINEAR PERTURBATIONS AND NON-GAUSSIANITY

Given that the two-point correlation function of the scalar perturbations carries no information about the non-Gaussian feature of the primordial perturbation, to study this feature it is necessary to calculate the three-point correlation function in our setup. By expanding the action \( \ref{eq:1} \) up to the third order in the small perturbations, eliminating \( R \) by using the equation \( \ref{eq:18} \) and introducing the auxiliary parameter \( \chi \) with
\[ Y = \frac{2(\kappa^2 - 8H \dot{G})}{2\kappa^2 H - 24H^2 \dot{G}} + \frac{a^2 \chi}{\kappa^2 - 8H \dot{G}}, \tag{37} \]
and
\[ \partial^2 \chi = \mathcal{W}_s \mathcal{D}, \tag{38} \]
we get
\[ S_3 = \int dt d^3 x \left\{ \left[ \frac{3a^3}{\kappa^2 c_s^2} \left( \frac{1}{c_s^2} - 1 \right) \right] \left( \epsilon - 4\kappa^2 H \dot{G} \right) \right\} \mathcal{D} \mathcal{D}^2 \]
\[ + \left[ \frac{a}{\kappa^2} \left( \frac{1}{c_s^2} - 1 \right) \right] \left( \epsilon - 4\kappa^2 H \dot{G} \right) \right\} \mathcal{D} \partial \mathcal{D}^2 \]
\[ + \left[ \frac{a^3}{\kappa^2} \left( \frac{1}{c_s^2} \right) \right] \left( \frac{1}{c_s^2} - 1 \right) \left( \epsilon - 4\kappa^2 H \dot{G} \right) \right\} \mathcal{D}^3 \]
\[ - \left[ \frac{a^3}{c_s^2} \right] \left( \epsilon - 4\kappa^2 H \dot{G} \right) \right\} \mathcal{D} \partial \mathcal{D} \partial \mathcal{D} \right\}. \tag{39} \]
The 3-point correlators in the interaction picture are obtained as
\[ \left( \mathcal{D}(k_1) \mathcal{D}(k_2) \mathcal{D}(k_3) \right) = \left( 2\pi \right)^3 \delta^3(k_1 + k_2 + k_3) \mathcal{B}_D(k_1, k_2, k_3), \tag{40} \]
where
\[ \mathcal{B}_D(k_1, k_2, k_3) = \left( 2\pi \right)^3 \frac{A^2}{\prod_{i=1}^{3} k_i^3} \mathcal{N}_D(k_1, k_2, k_3). \tag{41} \]
The parameter \( \mathcal{N}_D \) is defined by the following expression
\[ \mathcal{N}_D = \frac{3}{4} \left( 1 - \frac{1}{c_s^2} \right) \left( 2 \sum_{i>j} k_i^2 k_j^2 \right) \frac{1}{(k_1 + k_2 + k_3)^2} \]
\[ - \frac{1}{4} \left( 1 - \frac{1}{c_s^2} \right) \left( 2 \sum_{i>j} k_i^2 k_j^2 \right) \frac{1}{(k_1 + k_2 + k_3)^2} \]
\[ + \frac{1}{2} \sum_{i} k_i^4 \right\} \left[ \frac{1}{c_s^2} - 1 \right] \left( \frac{(k_1 k_2 k_3)^2}{(k_1 + k_2 + k_3)^3} \right), \tag{42} \]
by which we can define the so-called “nonlinearity parameter”, measuring the amplitude of the non-Gaussianity of the primordial perturbations, as follows
\[ f_{NL} = \frac{10}{3} \frac{\mathcal{N}_D}{\sum_{i=1}^{3} k_i^3}. \tag{43} \]
FIG. 1. Tensor-to-scalar ratio versus the scalar spectral index for a GB-α-attractor model in the background of Planck2015 TT, TE, EE+lowP data. As figure shows, when $\alpha \to 0$ all trajectories converge to the point shown by the red star, corresponding to $r \to 0$ and $n_s \to 0.967$. Also, when $\alpha \to \infty$, the trajectories tend to the points corresponding to the ones with $V \sim \phi^2$ and $G \sim G_0 \phi^2$. We have chosen $N = 60$ for the number of e-folds of inflation after the relevant CMB scales cross the horizon.

TABLE I. The ranges of $G_0$ for which the values of the scalar spectral index and the tensor-to-scalar ratio are compatible with Planck2015 observational data.

| $\alpha$       | $G_0$ range          |
|----------------|----------------------|
| $\alpha = 1$  | $G_0 < 4.1 \times 10^{-2}$ |
| $\alpha = 10^2$ | $G_0 < 3.91 \times 10^{-2}$ |
| $\alpha = 10^3$ | $G_0 < 2.83 \times 10^{-2}$ |
| $\alpha \to \infty$ | $0.217 < G_0 < 0.41$ |

The different shapes of the non-Gaussianity are obtained depending on the different values of the three momenta $k_1$, $k_2$, and $k_3$ (for details in this subject see [83, 86, 87, 88]). In a simple single field inflation model, where the non-Gaussianity is produced outside the horizon, the non-Gaussian feature is described with "local" type (where $k_3 \ll k_1 \approx k_2$) [13, 86, 89]. However, there are some models (such as the DBI, k-inflation and also higher derivative models) where the non-Gaussianity is created at horizon crossing during inflation. In these models the bispectrum has a maximal signal when all three wavelengths are equal to the horizon size ($k_1 = k_2 = k_3$) [22, 80]. In these models when any individual mode is far outside the horizon, the non-Gaussian feature is suppressed. In this regard, it is useful to study the non-Gaussianity in the "equilateral" configuration in these models. So, in our GB-α-attractor model, we study the non-Gaussian feature in the equilateral configuration. In this limit we have

$$N_D^{\text{equil}} = \frac{17}{72} k^3 \left( 1 - \frac{1}{c_s^2} \right),$$ 

leading to

$$f_{NL}^{\text{equil}} = \frac{85}{324} \left( 1 - \frac{1}{c_s^2} \right).$$ 

By using this equation and equations (1), (3) and (45), we can perform some numerical analysis on the model's parameters. The results are shown in figure 2. As this figure shows, in this GB-α-attractor model there is a set of model parameters which makes $f_{NL}^{\text{equil}}$ large enough to be detected in future surveys.

V. REHEATING

In this section, we are going to study the reheating era after the inflation. This process, which during it the universe reheats, helps us to obtain some additional constraints on the model’s parameters. In this regard, we follow the strategy used in Refs. [72, 73, 74, 75, 76] and find some expressions for $N_{rh}$ and $T_{rh}$. We can write these parameters in terms of the scalar spectral index $n_s$ by which we obtain constraints on $N_{rh}$ and $T_{rh}$. The number of e-folds between horizon crossing of the physi...
corresponding to $f_{NL} \rightarrow 0$ and $c_s^2 \rightarrow 1$. Also, when $\alpha \rightarrow \infty$, the trajectories tend to the points corresponding to the ones with $V \sim \frac{\phi^2}{2}$ and $G \sim G_0 \frac{\phi^2}{2}$. We have chosen $N = 60$ for the number of e-folds of inflation after the relevant CMB scales cross the horizon.

FIG. 2. Equilateral configuration of the non-Gaussianity versus the sound speed of the perturbations for a GB-$\alpha$-attractor model in the background of Planck2015 TTT, EEE, TTE and EET data. When $\alpha \rightarrow 0$ all trajectories converge to the point corresponding to $f_{NL} \rightarrow 0$ and $c_s^2 \rightarrow 1$. Also, when $\alpha \rightarrow \infty$, the trajectories tend to the points corresponding to the ones with $V \sim \frac{\phi^2}{2}$ and $G \sim G_0 \frac{\phi^2}{2}$. We have chosen $N = 60$ for the number of e-folds of inflation after the relevant CMB scales cross the horizon.

ical scales and the end of inflation phase is obtained as (see equation [13])

$$N_{hc} = \ln \left( \frac{a_f}{a_{hc}} \right) \simeq \int_{\phi_{hc}}^{\phi_f} \frac{-3H^2}{V' + 24H^4G'^2} d\phi,$$

with $a_f$ as the scale factor at the end of inflation and $a_{hc}$ as the scale factor at the horizon crossing. Assuming that during the reheating era, an energy density with the effective equation of state $\omega_{eff}$ dominates the universe, we can write $\rho \sim a^{-3(1+\omega_{eff})}$. We note that with a massive inflaton, $\omega_{eff}$ can be $-1$ (corresponding to the potential domination) and $+1$ (corresponding to the kinetic domination). However, at the end of the inflation epoch the value of the effective equation of state parameter is equal to $-\frac{3}{2}$. On the other hand, at the beginning of the radiation dominated universe, the value of this parameter is $\frac{1}{3}$. In this regard, it seems that the effective equation of state parameter during the reheating epoch is in the range $-\frac{3}{2}$ to $\frac{1}{3}$. By considering the massive inflaton, the frequency of the field’s oscillations is very larger than the rate of the expansion at the initial epoch of the reheating. So, the averaged effective pressure can be neglected and with a good approximation, the effective equation of state parameter can be considered to be zero. This means that at the beginning of the reheating epoch the value of $\omega_{eff}$ is effectively equal to the equation of state parameter of the matter. Then, by oscillation of the inflaton field and decaying to other particles, the value of the effective equation of state parameter increases with time and reaches $\frac{1}{3}$ at the beginning of the radiation domination era. The authors of Ref. [91] have indicated that in the chaotic inflation with a quartic interaction, the values of $\omega_{eff}$ are in the range between 0 and 0.2–0.3. Regarding that a constant effective equation of state parameter is allowed by a variety of reheating scenarios, we consider this parameter as a constant in our analysis. Now, we obtain the number of e-folds of the reheating era as

$$N_{rh} = \ln \left( \frac{a_{rh}}{a_f} \right) = -\frac{1}{3(1+\omega_{eff})} \ln \left( \frac{\rho_{rh}}{\rho_f} \right),$$

where the subscript $rh$ demonstrates the parameter in the reheating epoch. On the other hand, the energy density during inflation is given by

$$\rho = \left[ 1 + \frac{\epsilon^3}{3} - \frac{16\sqrt{2}k^3}{9} \sqrt{\epsilon} G'V + \left( \frac{32}{27} V^5 - \frac{64}{9} V^2 \right) \kappa^4 G'^2 \right] V.$$  

(48)

Now, if we show the value of $k$ at horizon crossing by $k_{hc}$, then we have

$$0 = \ln \left( \frac{k_{hc}}{a_{hc}H_{hc}} \right) = \ln \left( \frac{a_f}{a_{hc}} \frac{a_{rh}}{a_{hc}} \frac{a_0}{a_{rh}} \frac{k_{hc}}{H_{hc}} \right).$$

(49)

where $a_0$ is the current value of the scale factor. Equation 49 together with equations 46 and 47 lead to

$$N_{hc} + N_{rh} + \ln \left( \frac{a_0}{a_{rh}} \right) + \ln \left( \frac{k_{hc}}{a_0H_{hc}} \right) = 0.$$  

(50)

At this point we should find an expression for $\frac{a_0}{a_{rh}}$. The energy density of the reheating epoch has the following relation with the temperature of this epoch $T_{rh}$

$$\rho_{rh} = \frac{\pi^2 g_{rh}}{30} \frac{T^4_{rh}}{T_0^4},$$

(51)

where $g_{rh}$ demonstrates the effective number of relativistic species at the reheating. Also, the conservation of the entropy gives $T_{rh}$

$$\frac{a_0}{a_{rh}} = \left( \frac{43}{11g_{rh}} \right)^{\frac{1}{4}} \frac{T_{rh}}{T_0}.$$  

(52)
with $T_0$ being the current temperature of the universe. By using equation (51) and (52) we reach

$$\frac{a_0}{a_{rh}} = \left(\frac{43}{11g_{rh}}\right)^{-\frac{1}{3}} T_0^{-1} \left(\frac{\pi^2 g_{rh}}{30}\right)^{-\frac{1}{3}}. \quad (53)$$

To write $\rho_{rh}$ in terms of $N_{rh}$ from equation (47), we should first find $\rho_f$. The energy density at the end of inflation can be obtained by using equation (48) at the end of inflation (corresponding to $\epsilon = 1$). So

$$\rho_f = \left[ \frac{4}{3} - \frac{16\sqrt{2}\kappa^3}{9} G^\prime_f V_f + \left(\frac{32}{27} V_f^5 - \frac{64}{9} V_f^2\right) \kappa^6 G_f^2 \right] V_f. \quad (54)$$

From equations (47) and (53) we have

$$\rho_{rh} = \left[ \frac{4}{3} - \frac{16\sqrt{2}\kappa^3}{9} G^\prime_f V_f + \left(\frac{32}{27} V_f^5 - \frac{64}{9} V_f^2\right) \kappa^6 G_f^2 \right] V_f$$

$$\times \exp \left[ -3N_{rh}(1 + \omega_{eff}) \right]. \quad (55)$$

By using equation (53) and (55) we get

$$\ln \left(\frac{a_0}{a_{rh}}\right) = -\frac{1}{3} \ln \left(\frac{43}{11g_{rh}}\right) - \frac{1}{4} \ln \left(\frac{\pi^2 g_{rh}}{30}\right) - \ln T_0 +$$

$$\frac{1}{4} \ln \left[ \frac{4}{3} - \frac{16\sqrt{2}\kappa^3}{9} G^\prime_f V_f + \left(\frac{32}{27} V_f^5 - \frac{64}{9} V_f^2\right) \kappa^6 G_f^2 \right] V_f$$

$$\times \exp \left[ -\frac{3}{4} N_{rh}(1 + \omega_{eff}) \right]. \quad (56)$$

The value of the Hubble parameter at the horizon crossing ($H_{hc}$) is obtained from equation (24) as follows

$$H_{hc} = \sqrt{8A_0 W_0 \pi c_s^2}. \quad (57)$$

Now, using equations (50), (56) and (57) gives

$$N_{rh} = \frac{4}{1 - 3\omega_{eff}} \left[ -N_{hc} - \ln \left(\frac{k_{hc}}{a_0 T_0}\right) - \frac{1}{4} \ln \left(\frac{40}{\pi^2 g_{rh}}\right)$$

$$- \frac{1}{3} \ln \left(\frac{11g_{rh}}{43}\right) + \frac{1}{2} \ln \left(8\pi^2 A_0 W_0 c_s^2\right) - \frac{1}{4} \ln \left(\frac{4}{3}\right)$$

$$- \frac{16\sqrt{2}\kappa^3}{9} G^\prime_f V_f + \left(\frac{32}{27} V_f^5 - \frac{64}{9} V_f^2\right) \kappa^6 G_f^2 \right] V_f \right]. \quad (58)$$

Also, from equations (47), (52) and (54) we have

$$T_{rh} = \left(\frac{30}{\pi^2 g_{rh}}\right)^{\frac{1}{4}}$$

$$\times \left[ \frac{4}{3} - \frac{16\sqrt{2}\kappa^3}{9} G^\prime_f V_f + \left(\frac{32}{27} V_f^5 - \frac{64}{9} V_f^2\right) \kappa^6 G_f^2 \right]^{\frac{1}{4}} V_f$$

$$\times \exp \left[ -\frac{3}{4} N_{rh}(1 + \omega_{eff}) \right]. \quad (59)$$

By substituting equations (1) and (3) in equations (58) and (59), we obtain the e-folds number and temperature during reheating in terms of the model’s parameters, especially $\phi_{hc}$. On the other hand, $\phi_{hc}$ is related to the scalar spectral index, as is seen from equation (53) at the horizon crossing. So, it is possible to study the number of e-folds and temperature at the reheating versus the scalar spectral index. In our numerical analysis, we have adopted some sample values of $\omega_{eff}$ which are in the range between $-\frac{1}{3}$ and 1. We note that the values $-\frac{1}{3}$ and 1 are the very conservative allowed values. The results are shown in figures 3 and 4. Note that in plotting the figures and finding the constraints, we have used the value of the scalar spectral index as $n_s = 0.9652 \pm 0.0047$ from Planck2015 TT,TE,EE+lowP. As the figures show, for each class of the model’s parameters values, there is a maximum value for the number of e-folds which is consistent with observational data. In table II we have summarized the results for some values of the model’s parameters. Also, the figures show that there is a point that all curves in the plots originate from it. This point is corresponding to $N_{rh} = 0$ as instantaneous reheating. Note that, for GB coupling constant of the order of $10^{-2}$ or more and $\alpha$ of the order of 10, the instantaneous reheating is disfavored by Planck2015 observational data. We have found that, for all values of the parameter $\alpha$, as the GB coupling constant $G_0$ becomes smaller, the larger values of the e-folds number are more viable. Similarly, we can study the temperature in reheating phase numerically. For each class of the model’s parameter values there is a minimum temperature which is consistent with observational data (see table III). It is turned out that, for a given value of $\alpha$, as $G_0$ becomes smaller, a wider range of the reheating temperature is viable.

It should be noticed that, for $\omega_{eff} = \frac{1}{3}$ we can’t use equation (58). Actually even we repeat the analysis from equation (47) with $\omega_{eff} = \frac{1}{3}$, we can’t find an explicit expression for number of e-folds and temperature. However, it is possible to almost predict the value of the scalar spectral index. The curve for $\omega_{eff} = \frac{1}{3}$ in the plots can be a vertical line that crosses the instantaneous reheating point (see 73 [4]).

**VI. SUMMARY**

In this paper, we have considered a non-minimal Gauss-Bonnet inflation in the spirit of attractor scenario. Motivated by the $\alpha$-attractor models, we have adopted the potential and Gauss-Bonnet coupling function to be the same as the E-model potential. In this regard, in the small $\alpha$ limit there is an attractor point and in the large $\alpha$ limit we obtain $V \sim \phi^{2n}$ and $G \sim \phi^{2n}$. We have analyzed the perturbations in this setup and obtained the important perturbation parameters, such as the scalar spectral index and tensor-to-scalar ratio. To test the observational viability of this GB-$\alpha$-attractor model, we have studied $r$ and $n_s$ numerically and
TABLE II. The ranges of the number of e-folds parameter at reheating that are consistent with observational data.

| $\alpha$ | $G_0$ | $N_{rh} \leq$ | $N_{rh} \leq$ | $N_{rh} \leq$ |
|----------|-------|---------------|---------------|---------------|
| 0.1      | $8 \times 10^{-3}$ | 15             | 6             | 13            |
|          | $2 \times 10^{-2}$ | 24             | 10            | 17            |
|          | $8 \times 10^{-3}$ | 78             | 36            | 41            |
|          | $2 \times 10^{-2}$ | 50             | 80            | 68            |
|          | $4 \times 10^{-2}$ | 4              |               |               |

FIG. 3. Number of e-folds in reheating epoch versus the scalar spectral index for an E-model GB-$\alpha$-attractor inflation, with $n = 1$. The magenta region shows the values of the scalar spectral index released by Planck2015 data.

compared the results with Planck2015 data for $n = 1$. As expected, in small $\alpha$ limit the tensor-to-scalar ratio for all values of $G_0$ tends to zero and the scalar spectral index tends to 0.967. In the large $\alpha$ limit the model is corresponding to the model with $V \sim \phi^2$ and $\mathcal{G} \sim \phi^2$.

Our study shows that the GB-$\alpha$-attractor model in some domains of the model’s parameter space is consistent with observational data. The non-Gaussian feature of the primordial perturbations also, have been studied in this setup and the amplitude of the non-Gaussianity has
TABLE III. The ranges of the temperature at reheating which are consistent with observational data.

| Planck  | $\omega = -\frac{1}{3}$ | $\omega = 0$ | $\omega = \frac{1}{6}$ | $\omega = 0$ | $\omega = 1$ |
|---------|----------------|-------------|----------------|-------------|-------------|
| $\alpha = 0.1, G_0 = 8 \times 10^{-3}$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 13.3$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 10.67$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq -0.16$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 4.39$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 4.39$ |
| $\alpha = 0.1, G_0 = 2 \times 10^{-2}$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 13.65$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 13.29$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 5.71$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 1.18$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 2.95$ |
| $\alpha = 10, G_0 = 8 \times 10^{-3}$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 12.71$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 11.38$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 5.43$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 1.18$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \leq 13.94$ |
| $\alpha = 10, G_0 = 2 \times 10^{-2}$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 12.71$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 11.38$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 5.43$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \geq 1.18$ | $\log_{10} \left( \frac{T_{rh}}{\text{GeV}} \right) \leq 13.94$ |

FIG. 4. Temperature in reheating epoch versus the scalar spectral index for an E-model GB-$\alpha$-attractor inflation with $n = 1$. The light green region shows the temperatures below the electroweak scale, $T < 100$ GeV, and the dark green region demonstrates the temperatures below the big bang nucleosynthesis scale, $T < 10 MeV$. 

been obtained. We have obtained $f_{NL}$ in the equilateral configuration of the non-Gaussianity characterized by $k_1 = k_2 = k_3$. By studying the behavior of $f_{NL}^{\text{equil}}$ versus the sound speed, we have found that in small limit of $\alpha$ the amplitude of non-Gaussianity tends to zero corresponding to $c_s \to 1$. However, in the large limit of $\alpha$, depending on the GB coupling constant $G_0$, the sound speed can be small and $f_{NL}^{\text{equil}}$ can be large. So, in this GB-$\alpha$-attractor model there is a set of model parameters which makes $f_{NL}^{\text{equil}}$ large enough to be detected in future surveys. The reheating era
after the inflation epoch has been explored in this GB-$\alpha$-attractor model to obtain some more constraints on the model’s parameters. By assuming the domination of an energy density with an effective equation of state and considering its relation with the temperature, we have obtained the number of e-folds and temperature of the reheating in terms of the model’s parameters. After that, regarding to the relation of the scalar field at horizon crossing with the scalar spectral index, we have performed the numerical analysis on the behavior of $N_{rh}$ and $T_{rh}$ versus $n_s$. The results have been compared to Planck2015 constraints on the scalar spectral index. In this regard, we have obtained some constraints on the e-folds number and temperature of the reheating which are corresponding to the observationally viable values of the scalar spectral index. As an important result, we found that, for all values of the parameter $\alpha$, as the GB coupling constant reduces, the larger values of the e-folds number are more viable observationally.

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