

## Composite inflation confronts BICEP2 and PLANCK

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# Composite inflation confronts BICEP2 and PLANCK

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Received April 5, 2014

Revised May 20, 2014

Accepted May 31, 2014

Published June 19, 2014

**Abstract.** We examine observational constraints on single-field inflation in which the inflaton is a composite field stemming from a four-dimensional strongly interacting field theory. We confront the predictions with the Planck and very recent BICEP2 data. In the large non-minimal coupling regions, we discover for the minimal composite inflationary model that the predictions lie well inside the joint 68% CL for the Planck data, but is in tension with the recent BICEP2 observations. In the case of the glueball inflationary model, the predictions satisfy the Planck results. However, this model can produce a large tensor-to-scalar ratio consistent with the recent BICEP2 observations if the number of e-foldings is slightly smaller than the range commonly used. For a super Yang-Mills paradigm, we discover that the predictions satisfy the Planck data, and surprisingly a large tensor-to-scalar ratio consistent with the BICEP2 results can also be produced for an acceptable range of the number of e-foldings and of the confining scale. In the small non-minimal coupling regions, all of the models can satisfy the BICEP2 results. However, the predictions of the glueball and superglueball inflationary models cannot satisfy the observational bound on the amplitude of the curvature perturbation launched by Planck, and the techni-inflaton self-coupling in the minimal composite inflationary model is constrained to be extremely small.

**Keywords:** inflation, alternatives to inflation, particle physics - cosmology connection, modified gravity

**ArXiv ePrint:** [1307.2880](https://arxiv.org/abs/1307.2880)



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**1 Introduction**

It was widely expected that there was a period of accelerating expansion in the very early universe. Such period is traditionally known as inflation. The inflationary paradigm [1–5] tends to solve important issues, e.g. the magnetic monopoles, the flatness, and the horizon problems, plagued the standard big bang theory and successfully describes the generation and evolution of the observed large-scale structures of the universe. The inflationary scenario is formulated so far by the introduction of (elementary) scalar fields (called inflaton) with a nearly flat potential (see, e.g. [11–16]). However, the fundamental scalar field sector in field theories is plagued by the so-called hierarchy problem. Commonly, this means that quantum corrections generate unprotected quadratic divergences which must be fine-tuned away if the models must be true till the Planck energy scale. Similarly the inflaton, the field needed to initiate a period of rapid expansion of our Universe, suffers from the same kind of untamed quantum corrections. Therefore, finding its graceful exit is one of the great campaigns. Among other things, it seems like we have at least two compelling paradigms to avoid/solve such problem, namely Technicolor (TC) and Supersymmetry (SUSY).

All the extensions of the SM that address the naturalness problem tend to introduce a new symmetry in which the scalar fields are linked to fermionic ones. In the case of TC the symmetry that protects the scalar mass is the chiral symmetry of the constituent fermions. In the same manner, a natural hope for employing the composite fields to cosmology is also to solve such problem, as well. Therefore, it would be of great interest to imagine natural models underlying the cosmic inflation in which the inflaton need not be an elementary degree of freedom. Recent investigations show that it is possible to construct models in which the inflaton emerges as a composite state of a four-dimensional strongly coupled theory [18–20]. These types of models have already been stamped to be composite inflation. There were other models of super or holographic composite inflation [36–43].

Based on the composite inspiration, we start with the basic assumption that the inflaton needs not be an elementary degree of freedom emerging from a strongly interacting field theory. The authors of [18–20] have already examined the paradigms in which inflation is driven by a four-dimensional strongly coupled dynamic with a nonminimal coupling to gravity. They also demonstrated that it is possible to achieve successful inflation with the confining scale of the underlying gauge theory naturally of the order of the grand unified energy scale. Moreover, they showed that, all models presented in this present work nicely respect tree-level unitarity for the scattering of the inflaton field all the way to the Planck scale.

Practically, all speculative ideas concerning physics of very early universe can be falsified by using the observation of the large scale structure. In particular, the temperature fluctuations observed in the Cosmic Microwave Background (CMB) is basically regarded as providing a clear window to probe the inflationary cosmology. To testify all inflationary models, we need observables including: (i) the scalar spectral index  $n_s$ , (ii) the amplitude of the power spectrum for the curvature perturbations, (iii) tensor-to-scalar ratio  $r$ , (iv) the non-gaussianity parameter  $f_{\text{NL}}$ .

Yet, other relevant parameters are the running of scalar spectral index  $\alpha \equiv dn_s/d\ln k$  and the spectral index for tensor perturbations  $n_T$ . Most recently, the Planck satellite data showed that the spectral index  $n_s$  of curvature perturbations is constrained to be  $n_s = 0.9603 \pm 0.0073$  (68% CL) and ruled out the exact scale-invariance ( $n_s = 1$ ) at more than  $5\sigma$  confident level (CL), whilst the amplitude of the power spectrum for the curvature perturbations  $|\zeta|^2$  is bounded to be  $\mathcal{A}_s = 3.089_{-0.027}^{+0.024}$  (68% CL) [21] with  $\mathcal{A}_s \equiv \ln(|\zeta|^2 \times 10^{10})$ . Having constrained by Planck, the tensor-to-scalar ratio  $r$  is bounded to be  $r < 0.11$  (95%CL). Surprisingly, the recent BICEP2 data renders the bound on  $r$  to be  $r = 0.20_{-0.05}^{+0.07}$  with  $r = 0$  disfavored at  $7.0\sigma$  CL [22]. The attempts to explain the seed perturbations from inflation were originally studied by many authors [6–9]. As for gravitational waves generated during inflation, it was earlier calculated in [10].

In the context of composite inflation, the previous publications demonstrated that it was significantly possible to achieve successful inflation with the amplitude of scalar perturbations compatible with the CMB constraint. However, the previous examinations can be observationally tested if we compute the observables, e.g. the tensor-to-scalar ratio and the spectral index, predicted by the theory on the subjects and constrain the resulting predictions by using the results from the Planck data and very recent BICEP2 experiments.

The paper is organized as follows: in section 2, we first spell out the setup for a generic model of composite paradigm. We then derive equations of motion and useful expressions. In section 3, we derive  $n_s$ ,  $r$  and  $\mathcal{A}_s$  for composite paradigms. In section 4, we examine  $n_s$ ,  $r$  and  $\mathcal{A}_s$ , and the range of the model parameters in which these quantities satisfy the observational bound is evaluated. Using the Planck and very recent BICEP2 observations of the  $(n_s - r)$  plane, we also confront our results with such observations. Finally, the conclusions are given in section 5.

## 2 Composite formulations and background evolutions

Recently, it has already been shown that cosmic inflation can be driven by four-dimensional strongly interacting theories non-minimally coupled to gravity [18–20]. The general action for composite inflation in the Jordan frame takes the form for scalar-tensor theory of gravity as<sup>1</sup>

$$\mathcal{S}_{\text{CI,J}} = \int d^4x \sqrt{-g} \left[ \frac{M_{\text{P}}^2}{2} F(\Phi) R - \frac{1}{2} G(\Phi) g^{\mu\nu} \partial_\mu \Phi \partial_\nu \Phi - V(\Phi) \right]. \quad (2.1)$$

<sup>1</sup>We used the signature of the matrix as  $(-, +, +, +)$  throughout the paper.

The functions  $F(\Phi)$  and  $G(\Phi)$  in this action are defined as

$$F(\Phi) = 1 + \frac{\xi}{M_{\text{P}}^2} \Phi^{\frac{2}{D}} \quad \text{and} \quad G(\Phi) = \frac{1}{D^2} G_0 \Phi^{\frac{2-2D}{D}}, \quad (2.2)$$

where  $D$  is the mass dimension of the composite field  $\Phi$ ,  $G_0$  is a constant and  $1/D^2$  is introduced for later simplification. In our setup, we write the potential in the following form:

$$V(\Phi) = \Phi^{4/D} f(\Phi) \quad \text{with} \quad \Phi \equiv \varphi^D, \quad (2.3)$$

where the field  $\varphi$  possesses unity canonical dimension and  $f(\Phi)$  is a general function of the field  $\Phi$  concretely implemented below. The non-minimal coupling to gravity is signified by the dimensionless coupling  $\xi$ . Here, we write the general action for the composite inflation in the form of scalar-tensor theory of gravity in which the inflaton non-minimally couples to gravity. At the moment, the non-minimal term  $\xi \Phi^{2/D} R/M_{\text{P}}^2$  has purely phenomenological origin. The reason resides from the fact that one wants to relax the unacceptable large amplitude of primordial power spectrum generated if one takes  $\xi = 0$  or smaller. This non-minimal coupling of scalar fields to gravity was pioneered in several earlier [44–49] and recent works [50] where a similar phenomenological large value of  $\xi$  was needed. However, it is very interesting to further study a potential origin of such a large coupling.

According to the above action, the Friedmann equation and the evolution equations for the background field are respectively given by

$$3M_{\text{P}}^2 F H^2 + 3M_{\text{P}}^2 \dot{F} H = 3M_{\text{P}}^2 H^2 F (1 + 2\mathcal{F}_t) = \frac{1}{2} G \dot{\Phi}^2 + V(\Phi), \quad (2.4)$$

$$3M_{\text{P}}^2 F H^2 + 2M_{\text{P}}^2 \dot{F} H + 2M_{\text{P}}^2 F \dot{H} + M_{\text{P}}^2 \ddot{F} = -\frac{1}{2} G \dot{\Phi}^2 + V(\Phi), \quad (2.5)$$

$$G \ddot{\Phi} + 3HG\dot{\Phi} + \frac{1}{2} G_{\Phi} \dot{\Phi}^2 + V_{\Phi} = 3M_{\text{P}}^2 F_{\Phi} (\dot{H} + 2H^2), \quad (2.6)$$

where  $\mathcal{F}_t = \dot{F}/(2HF)$ ,  $H$  is the Hubble parameter, subscripts “ $\Phi$ ” denote derivative with respect to  $\Phi$ , and the dot represents derivative with respect to time,  $t$ . In the following calculations, we will set  $M_{\text{P}}^2 = 1$ . In order to derive the observables, it is common to apply the standard slow-roll approximations such that

$$|\ddot{\Phi}/\dot{\Phi}| \ll H, \quad |\dot{\Phi}/\Phi| \ll H \quad \text{and} \quad |G\dot{\Phi}^2/2| \ll V(\Phi). \quad (2.7)$$

With these conditions, the equations given in eqs. (2.4)–(2.7) become

$$\epsilon = \mathcal{F}_t - \frac{V_{\Phi} F}{V F_{\Phi}} \mathcal{F}_t \quad \text{and} \quad \Phi' = \frac{1}{\left(1 + \frac{3F_{\Phi}^2}{2FG}\right)} \left(2\frac{F_{\Phi}}{G} - \frac{V_{\Phi} F}{V G}\right). \quad (2.8)$$

Here the first relation can be derived by differentiating eq. (2.4) with respect to time and then applying the above slow-roll conditions and  $\Phi' \equiv d\Phi/d \ln a = \dot{\Phi}/H$ . In order to keep our investigation more transparency, we will express hereafter the relevant parameters in term of the canonical-dimension one field  $\varphi$ . For more convenience, we expand the  $\epsilon$  parameters around the coupling  $\xi$ , and find for  $\xi \ll 1$

$$\epsilon \simeq \frac{(4 + \Theta\varphi)^2}{2G_0\varphi^2} + \frac{1}{2G_0} \left[ -8 + 2\Theta\varphi + \Theta^2\varphi^2 \right] \xi + \mathcal{O}(\xi^2), \quad (2.9)$$

and for  $\xi \gg 1$

$$\epsilon \simeq \frac{1}{12} [2\Theta\varphi + \Theta^2\varphi^2] + \frac{1}{72\xi\varphi^2} [48 + 60\Theta\varphi + 12\Theta^2\varphi^2 - 2G_0\Theta\varphi^3 - G_0\Theta^2\varphi^4] + \mathcal{O}(1/\xi^2), \quad (2.10)$$

with  $\Theta \equiv (1/f(\varphi))(\partial f(\varphi)/\partial\varphi)$ . We can also do the same exercise to expand  $\varphi'$  around a small  $\xi$  such that

$$\varphi' \simeq -\frac{1}{G_0\varphi} [4 + \Theta\varphi] - \frac{\Theta\varphi^2}{G_0}\xi + \mathcal{O}(\xi^2), \quad (2.11)$$

and for  $\xi \gg 1$

$$\varphi' \simeq -\frac{1}{6}\Theta\varphi^2 + \frac{\varphi^{-1}}{36\xi} [-24 - 12\Theta\varphi + G_0\Theta\varphi^3] + \mathcal{O}(1/\xi^2). \quad (2.12)$$

Having computed the field  $\varphi$  at the end of inflation  $\varphi_e$  by using the condition  $\epsilon(\varphi_e) = 1$ , one can determine the number of e-foldings via

$$\mathcal{N}(\varphi) = \ln \frac{a_e}{a} = \int_a^{a_e} \frac{1}{\tilde{a}} d\tilde{a} = \int_t^{t_e} H d\tilde{t} = \int_\varphi^{\varphi_e} \frac{H}{\dot{\tilde{\varphi}}} d\tilde{\varphi} = \int_\varphi^{\varphi_e} \frac{1}{\tilde{\varphi}'} d\tilde{\varphi}, \quad (2.13)$$

where the subscript “e” denotes the evaluation at the end of inflation and  $\varphi'$  is given by eq. (2.8). At the observable perturbation exits the horizon, we can evaluate the field  $\varphi$  once the number of e-foldings  $\mathcal{N}$  is specified. Determining the value of  $\varphi$  and  $\varphi'$  when the perturbations exit the horizon allows us to compute the spectral index and power spectrum amplitude in terms of the number of e-foldings.

### 3 Power spectra and spectral index

In order to obtain the power spectra for our models, it is (tricky) convenient to use the results in Einstein frame. In order to transform the action in the Jordan frame into the Einstein one, we take the following transformation,

$$\tilde{g}_{\mu\nu} = F(\Phi) g_{\mu\nu}. \quad (3.1)$$

Regarding to the above (conformal) implementation, the action in eq. (2.1) can be written in Einstein frame as

$$\mathcal{S}_{\text{Cl,E}} = \int d^4x \sqrt{-\tilde{g}} \left[ \frac{M_{\text{P}}^2}{2} \tilde{R} - \frac{1}{2} \partial_\mu \tilde{\Phi} \partial^\mu \tilde{\Phi} - U(\tilde{\Phi}) \right], \quad (3.2)$$

where  $\tilde{g}$  and  $\tilde{R}$  are computed from  $\tilde{g}_{\mu\nu}$ , “tildes” represent the quantities in the Einstein frame, and

$$\frac{\partial\Phi}{\partial\tilde{\Phi}} = \frac{F}{\sqrt{GF + 3F_\Phi^2/2}}, \quad \text{and} \quad U(\tilde{\Phi}) = \frac{V(\Phi)}{F^2(\Phi)} \Big|_{\Phi=\tilde{\Phi}(\tilde{\Phi})}. \quad (3.3)$$

Using the expression for the slow-roll parameter,  $\tilde{\epsilon}$ , in the Einstein frame such that

$$\tilde{\epsilon} = \frac{1}{2} \left( \frac{1}{U} \frac{\partial U}{\partial \tilde{\Phi}} \right)^2, \quad (3.4)$$

one can simply show that

$$\tilde{\epsilon} = \frac{1}{2} \left( \frac{F^2}{V} \frac{\partial \Phi}{\partial \tilde{\Phi}} \frac{\partial}{\partial \Phi} \left( \frac{V}{F^2} \right) \right)^2 = \epsilon + \mathcal{F}_t. \quad (3.5)$$

It is well known that the power spectrum for the scalar perturbation generated from inflaton field  $\tilde{\Phi}$  in the Einstein frame is given by

$$\mathcal{P}_\zeta \simeq \frac{U}{24\pi^2 \tilde{\epsilon}} \Big|_{k|\tau|=1}, \quad (3.6)$$

where the above expression is evaluated at the conformal time  $\tau$  when the perturbation with wavenumber  $k$  exits the horizon, and the tensor-to-scalar ratio is

$$r \simeq 16\tilde{\epsilon}. \quad (3.7)$$

Since the power spectra are frame-independent, we can use eq. (3.5) to write the power spectrum in eq. (3.6) and the tensor-to-scalar ratio in eq. (3.7) in terms of  $\epsilon$  as

$$\mathcal{P}_\zeta \simeq \frac{V}{24\pi^2 F(\epsilon + \mathcal{F}_t)} \Big|_{k|\tau|=1}, \quad (3.8)$$

$$r \simeq 16(\epsilon + \mathcal{F}_t). \quad (3.9)$$

The spectrum index for this power spectrum can be computed via

$$n_s = \frac{d \ln \mathcal{P}_\zeta}{d \ln k} + 1 \simeq \frac{1}{H} \frac{d \ln \mathcal{P}_\zeta}{dt} + 1 \simeq 1 - 2\epsilon - 2\mathcal{F}_t - \Phi' \frac{d \ln(\epsilon + \mathcal{F}_t)}{d\Phi}, \quad (3.10)$$

Using eq. (2.8), we get

$$\begin{aligned} n_s = 1 + \frac{1}{\varphi^2 (G_0 \varphi^2 \xi + G_0 + 6\varphi^2 \xi^2)^2} \times \\ \left[ 2 \left( (\varphi^3 \xi + \varphi)^2 \Theta_\varphi (G_0 \varphi^2 \xi + G_0 + 6\varphi^2 \xi^2) - 4(G_0 (2\varphi^4 \xi^2 + 5\varphi^2 \xi + 3) + 12\varphi^2 \xi^2 (\varphi^2 \xi + 2)) \right) \right. \\ \left. + 2\varphi \Theta (\varphi^2 \xi + 1) (G_0 (\varphi^4 \xi^2 - 3\varphi^2 \xi - 4) + 6\varphi^2 \xi^2 (\varphi^2 \xi - 5)) \right. \\ \left. + \Theta^2 \left( -(\varphi^3 \xi + \varphi)^2 (G_0 \varphi^2 \xi + G_0 + 6\varphi^2 \xi^2) \right) \right], \end{aligned} \quad (3.11)$$

where  $\Theta_\varphi = \partial \Theta / \partial \varphi$ . For our discussion below, we expand the spectrum index around  $\xi$  for both  $\xi \ll 1$  and  $\xi \gg 1$ . For the first case  $\xi \ll 1$ , we find

$$n_s \simeq 1 - \frac{(24 + 8\Theta\varphi + \Theta^2\varphi^2)}{G_0\varphi^2} + \frac{1}{G_0} [8 + 2\Theta\varphi - \Theta^2\varphi^2] \xi + \mathcal{O}(\xi^2), \quad (3.12)$$

and for  $\xi \gg 1$

$$\begin{aligned} n_s \simeq 1 + \frac{1}{6} \Theta\varphi (2 - \Theta\varphi) + \frac{1}{36\xi\varphi^2} \left( -96 - 48\Theta\varphi - 12\Theta^2\varphi^2 - 2\Theta G_0\varphi^3 + \Theta^2 G_0\varphi^4 \right) \\ + \mathcal{O}(1/\xi^2). \end{aligned} \quad (3.13)$$

The amplitude of the curvature perturbation can be directly read from the power spectrum and we find

$$|\zeta|^2 \simeq \frac{V}{24\pi^2 F^2 (\epsilon + \mathcal{F}_t)} \Bigg|_{c_s k|\tau|=1}, \quad (3.14)$$

Using eq. (2.8), the above equation becomes

$$|\zeta|^2 = \frac{\varphi^2 (G_0 \varphi^2 \xi + G_0 + 6\varphi^2 \xi^2) V(\varphi)}{12\pi^2 (\varphi^2 \xi + 1)^2 (\Theta (\varphi^3 \xi + \varphi) + 4)^2}. \quad (3.15)$$

The expansion of  $|\zeta|^2$  around  $\xi$  for both  $\xi \ll 1$  and  $\xi \gg 1$  respectively reads

$$|\zeta|^2 \simeq \frac{G_0 \varphi^6 f(\varphi)}{12\pi^2 (4 + \Theta \varphi)^2} - \frac{G_0 f(\varphi)}{12\pi^2 (4 + \Theta \varphi)^3} \left[ \varphi^8 (4 + 3\Theta \varphi) \right] \xi + \mathcal{O}(\xi^2), \quad (3.16)$$

and

$$|\zeta|^2 \simeq \frac{f(\varphi)}{2\pi^2 \Theta^2 \varphi^2 \xi^2} + \frac{\varphi^{-5}}{12\pi^2 \Theta^3 \xi^2} \left[ -48 - 24\Theta \varphi + G_0 \Theta \varphi^3 \right] f(\varphi) + \mathcal{O}(1/\xi^4), \quad (3.17)$$

We can write the tensor-to-scalar ratio in terms of functions  $F$  and  $G$  as

$$r = 16 \frac{(\Theta (\varphi^3 \xi + \varphi) + 4)^2}{2\varphi^2 (G_0 \varphi^2 \xi + G_0 + 6\varphi^2 \xi^2)}. \quad (3.18)$$

The  $\xi$ -expansion of the tensor-to-scalar ratio  $r$  reads for  $\xi \ll 1$

$$r \simeq \frac{8(4 + \Theta \varphi)^2}{G_0 \varphi^2} + \frac{8}{G_0} (-16 + \Theta^2 \varphi^2) \xi + \mathcal{O}(\xi^2). \quad (3.19)$$

and for  $\xi \gg 1$

$$r \simeq \frac{4\Theta^2 \varphi^2}{3} + \frac{2(48\Theta + 12\Theta^2 \varphi - G_0 \Theta^2 \varphi^3)}{9\xi \varphi} + \mathcal{O}(1/\xi^2). \quad (3.20)$$

In the next section, we will implement our setup to investigating the models of composite inflation.

## 4 Contact with observations

In this section, we examine observational constraints on single-field inflation in which the inflaton is a composite field stemming from a four-dimensional strongly interacting field theory non-minimally coupled to gravity using the Planck data and BICEP2 results.

### 4.1 Minimal composite inflation (MCI)

The authors of [18] recently demonstrated that it is possible to obtain a successful inflation in which the inflaton is a composite field stemming from a four-dimensional strongly interacting field theory. In this work, they engaged the simplest models of technicolor known as



the minimal walking technicolor (MWT) theory [27–30] with the standard (slow-roll) inflationary paradigm as a template for composite inflation. We discussed in the appendix A the underlying gauge theory of inflation and showed how the non-minimal coupling comes into play in the theory at the fundamental description. In this section, we further examine this model of composite inflation using Planck and BICEP2 data. In this model, the composite inflaton is the lightest state  $\varphi$ , and the remaining composite fields are massive. This provides a sensible possibility to consider the  $\varphi$  dynamics first. For model of inflation, we write  $\Phi \equiv \varphi$ , in which we couple non-minimally to gravity. The resulting action in the Jordan frame is given by [18]:

$$\mathcal{S}_{\text{MCI}} = \int d^4x \sqrt{-g} \left[ \frac{1 + \xi \varphi^2}{2} R - \frac{1}{2} g^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi - V_{\text{MCI}}(\varphi) \right], \quad (4.1)$$

where

$$V_{\text{MCI}}(\varphi) = -\frac{m^2}{2} \varphi^2 + \frac{\kappa}{4} \varphi^4, \quad (4.2)$$

in which  $\kappa$  is a self coupling and the inflaton mass is  $m_{\text{TI}}^2 = 2m^2$ . Since  $m_{\text{TI}}$  is order of the GeV energy scale,  $\kappa$  should be order of unity and  $\varphi$  during inflation is order of Planck mass, we neglect  $m_{\text{TI}}^2$  term in our calculation. For this model, we have

$$F(\varphi) = 1 + \xi \varphi^2 \quad \text{and} \quad G = G_0 = 1. \quad (4.3)$$

For this form of the potential, eqs. (3.11), (3.18), (3.15) and (2.13) yield the intuitive results,

$$n_s \simeq 1 - \frac{8(2\varphi^4 \xi^2 (6\xi + 1) + \varphi^2 \xi (24\xi + 5) + 3)}{(\varphi^3 \xi (6\xi + 1) + \varphi^2)^2}, \quad r \simeq \frac{128}{\varphi^4 \xi (6\xi + 1) + \varphi^2}, \quad (4.4)$$

$$|\zeta|^2 \simeq \frac{\kappa \varphi^6 (\varphi^2 \xi (6\xi + 1) + 1)}{768 (\pi \varphi^2 \xi + \pi)^2}, \quad \mathcal{N} \simeq \frac{1}{8} \left( (6\xi + 1) (\varphi^2 - \varphi_e^2) - 6 \log \left( \frac{\varphi^2 \xi + 1}{\varphi_e^2 \xi + 1} \right) \right). \quad (4.5)$$

For more transparency, we consider in the large  $\xi$  limit, and write  $n_s$ ,  $r$  and  $|\zeta|^2$  in terms of  $\mathcal{N}$  as

$$n_s \simeq 1 - \frac{8}{3\varphi^2 \xi} + \mathcal{O}(1/\xi^2) \simeq 1 - \frac{2}{\mathcal{N}}, \quad (4.6)$$

$$r \simeq \frac{64}{3\varphi^4 \xi^2} - \frac{32}{9\varphi^4 \xi^3} + \mathcal{O}(1/\xi^4) \simeq \frac{12}{\mathcal{N}^2}, \quad (4.7)$$

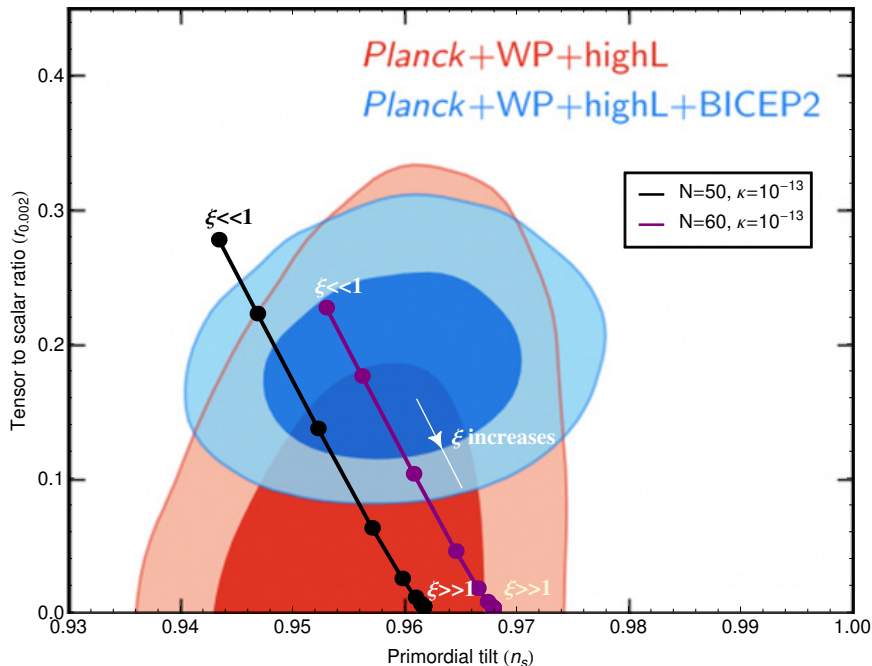
$$|\zeta|^2 \simeq \frac{\kappa \varphi^4}{128\pi^2} + \frac{\kappa(-12\varphi^2 + \varphi^4)}{768\pi^2 \xi} + \mathcal{O}(1/\xi^2) \simeq \frac{\kappa \mathcal{N}^2}{72\pi^2 \xi^2}. \quad (4.8)$$

In this case with  $\kappa \sim \mathcal{O}(1)$ ,  $n_s$ ,  $r$  and  $\mathcal{A}_s$  are well consistent with the Planck data up to  $2\sigma$  CL for  $\mathcal{N} = 60$  and  $4.7 \times 10^4 \lesssim \xi \lesssim 5.0 \times 10^4$ , for instance, illustrated in figure 2. However,  $\mathcal{A}_s$  does strongly depend on  $\mathcal{N}$ , and thus the coupling can be lowered (or raised) if  $\mathcal{N}$  changes. Unfortunately, with large  $\xi$ , the prediction  $r$  for this model is in tension with the recent BICEP2 results shown in figure 1. In the small  $\xi$  limit, we have

$$n_s \simeq \left( 1 - \frac{24}{\varphi^2} \right) + 8\xi + \mathcal{O}(\xi^2) \simeq 1 - \frac{3}{\mathcal{N}}, \quad (4.9)$$

$$r \simeq \frac{128}{\varphi^2} - 128\xi + \mathcal{O}(\xi^2) \simeq \frac{16}{\mathcal{N}}, \quad (4.10)$$

$$|\zeta|^2 \simeq \frac{\kappa \varphi^6}{768\pi^2} - \frac{\kappa \varphi^8 \xi}{768\pi^2} + \mathcal{O}(\xi^2) \simeq \frac{2\kappa \mathcal{N}^3}{3\pi^2}. \quad (4.11)$$



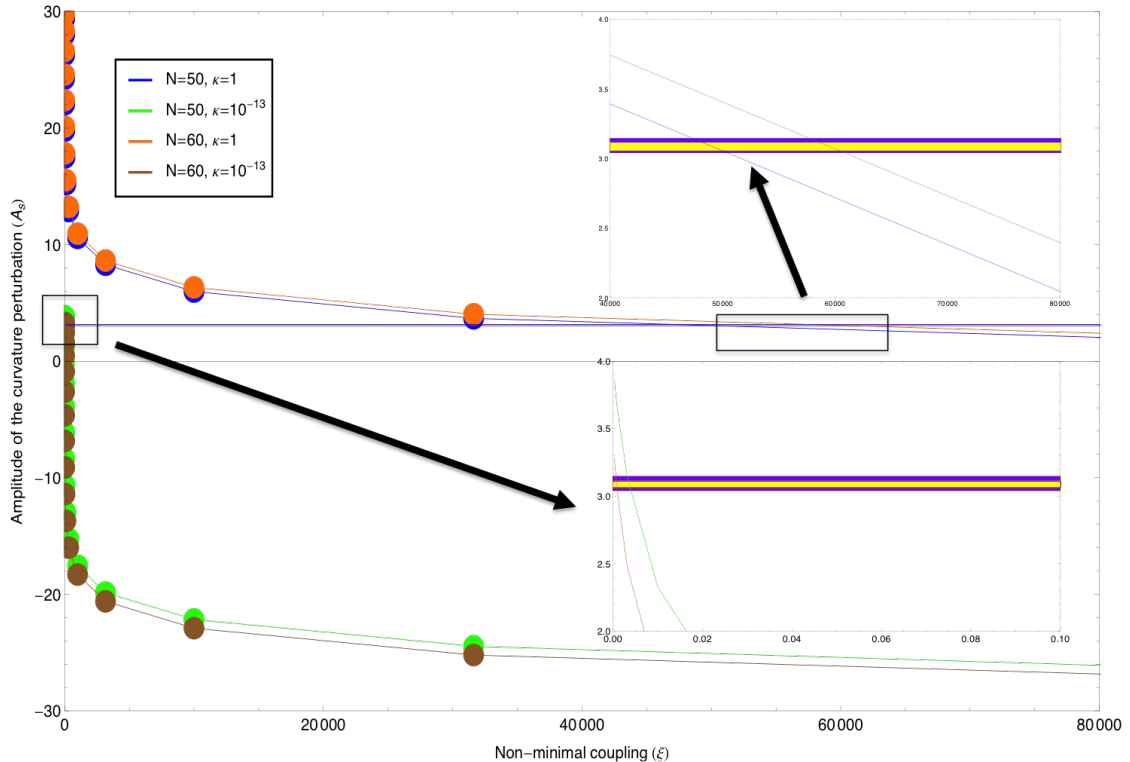
**Figure 1.** The contours display the resulting 68% and 95% confidence regions for the tensor-to-scalar ratio  $r$  and the scalar spectral index  $n_s$ . The red contours are for the Planck+WP+highL data combination, which for this model extension gives a 95% bound  $r < 0.26$  (Planck Collaboration XVI 2013 [21]). The rest represents the BICEP2 constraints on  $r$ . The plots show the numerical data resulting from the MCI model by varying the coupling  $\xi$  within a range of  $10^{-3} \lesssim \xi \lesssim 10^6$ .

In this situation, the prediction of  $r$  for this model can be consistent with the recent BICEP2 results shown in figure 1. In addition, we discover that the amplitude of the curvature perturbation  $\mathcal{A}_s$  can satisfy the Planck data at  $2\sigma$  CL for  $\mathcal{N} = 60$ ,  $\xi \sim 10^{-3}$  if  $\kappa \sim 10^{-13}$  illustrated in figure 2. However, in this case the prediction of such a very small  $\kappa$  is opposed to that from the underlying theory  $\kappa \sim \mathcal{O}(1)$ . We will further provide detailed discussions about this model for a small and large  $\xi$  in the last section.

## 4.2 Glueball inflation (GI)

In this section, we consider another viable model of composite inflation. The simplest, but intuitive, examples of strongly coupled theories are pure Yang-Mills theories featuring only gluonic-type fields. We discussed in the appendix B the underlying gauge theory of inflation and showed how the non-minimal coupling pops up in the theory at the fundamental description. The authors of [19] demonstrated that it is possible to achieve successful inflation where the inflaton emerges as the interpolating field describing the lightest glueball associated to a pure Yang-Mills theory. The original derivation of the low-energy effective Lagrangian can be found in [32–34]. It is worthy to note here that one can effectively solve the cosmological “hierarchy problem” in the scalar sector of the inflation which is not solved by Higgs inflation. For this model, we have

$$f(\Phi) = \frac{1}{2} \ln(\Phi/\Lambda^4), \quad (4.12)$$



**Figure 2.** The plot shows the relation between the amplitude of the power spectrum  $\mathcal{A}_s$  and the non-minimal coupling  $\xi$  within a range of  $10^{-3} \lesssim \xi \lesssim 10^6$  for  $\mathcal{N} = 50, 60$  predicted by the MCI model. The horizontal bands represent the  $1\sigma$  (yellow) and  $2\sigma$  (purple) CL for  $\mathcal{A}_s$  obtained from Planck.

so that the effective Lagrangian for the lightest glueball state, constrained by the Yang-Mills trace anomaly, non-minimally coupled to gravity in the Jordan frame reads

$$\mathcal{S}_{\text{GI}} = \int d^4x \sqrt{-g} \left[ \frac{1 + \xi \Phi^{1/2}}{2} R - \frac{1}{2} \Phi^{-3/2} g^{\mu\nu} \partial_\mu \Phi \partial_\nu \Phi - V_{\text{GI}}(\Phi) \right], \quad (4.13)$$

where

$$V_{\text{GI}}(\Phi) = \frac{\Phi}{2} \ln(\Phi/\Lambda^4), \quad (4.14)$$

Here the physical meaning of  $\Lambda$  underlying inflationary scenario has already been discussed in [19]. The effective potential given above is known in particle physics. It is the generating function for the trace anomaly of a generic purely gluonic Yang-Mills theory such that  $\Phi$  is a composite operator. It is convenient to introduce the field  $\varphi$  possessing unity canonical dimension and related to  $\Phi$  as follows:

$$\Phi = \varphi^4. \quad (4.15)$$

From the above assignment, the action then becomes

$$\mathcal{S}_{\text{GI}} = \int d^4x \sqrt{-g} \left[ \frac{1 + \xi \varphi^2}{2} R - \frac{1}{2} g^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi - V_{\text{GI}}(\varphi) \right], \quad (4.16)$$

where

$$V_{\text{GI}}(\varphi) = 2\varphi^4 \ln(\varphi/\Lambda), \quad (4.17)$$

which yields

$$F(\varphi) = 1 + \xi \varphi^2 \quad \text{and} \quad G_0 = 16. \quad (4.18)$$

Here the above quantities can be directly read from the action (4.16). For this form of the potential,  $F$  and  $G$ , we can use eq. (2.13) to compute the number of e-foldings in the  $\xi \gg 1$  and  $\xi \ll 1$  limits respectively as

$$\mathcal{N} \simeq 3 \left( \ln^2(\varphi/\Lambda) - \ln^2(\varphi_e/\Lambda) \right) + \mathcal{O}(1/\xi), \quad (4.19)$$

and

$$\mathcal{N} \simeq \int_{\varphi_e}^{\varphi} \frac{16\tilde{\varphi} \ln(\tilde{\varphi}/\Lambda) d\tilde{\varphi}}{4 \ln(\tilde{\varphi}/\Lambda) + 1} + \mathcal{O}(\xi) \simeq 2\varphi^2 - 2\varphi_e^2 + \mathcal{O}(\xi), \quad (4.20)$$

where in the calculation of  $\mathcal{N}$  for the small  $\xi$  case, we have supposed that  $\ln(\varphi/\Lambda) \gg 1$ . This is so since when  $\ln(\varphi/\Lambda) < 1$  we get  $\mathcal{N} < 0$  unless  $\varphi < \varphi_e$ . However, this approximation for  $\mathcal{N}$  at small  $\xi$  limit is valid if  $\Lambda \lesssim 1$ , since the  $\mathcal{O}(\xi)$  term in eq. (2.11) becomes significant when  $\Lambda > 1$ . In contrary, for  $\xi \gg 1$  case, we compute the expression for  $\mathcal{N}$  by supposing that  $\Lambda$  is not too small, otherwise we have to keep the next leading order terms in the expansion of  $\varphi'$  and  $\epsilon$  given in eqs. (2.10) and (2.12). In the case where  $\Lambda$  is not too small, it follows from eq. (2.10) that  $\ln(\varphi_e/\Lambda) \sim \mathcal{O}(1)$  at leading order. Moreover, we see from eq. (2.9) that  $\varphi_e \sim \mathcal{O}(1)$  when  $\Lambda < 1$ . Hence, for these cases, we neglect the  $\varphi_e$ -dependent terms in the above expression for  $\mathcal{N}$ . Using these expression for  $\mathcal{N}$ , we can write  $\varphi$  in terms of  $\mathcal{N}$ . Therefore we can use eqs. (3.13), (3.17) and (3.20) to write  $n_s$ ,  $r$  and  $|\zeta|^2$  in terms of  $\mathcal{N}$  for  $\xi \gg 1$  case as

$$n_s \simeq 1 - \frac{1}{2 \ln^2(\varphi/\Lambda)} + \mathcal{O}(1/\xi) \simeq 1 - \frac{3}{2\mathcal{N}} + \mathcal{O}(\xi), \quad (4.21)$$

$$r \simeq \frac{4}{3 \ln^2(\varphi/\Lambda)} + \mathcal{O}(1/\xi) \simeq \frac{4}{\mathcal{N}} + \mathcal{O}(\xi), \quad (4.22)$$

$$|\zeta|^2 \simeq \frac{\ln^3(\varphi/\Lambda)}{\pi^2 \xi^2} + \mathcal{O}(1/\xi^3) \simeq \frac{\mathcal{N}^{3/2}}{3\sqrt{3}\pi^2 \xi^2} + \mathcal{O}(1/\xi^3), \quad (4.23)$$

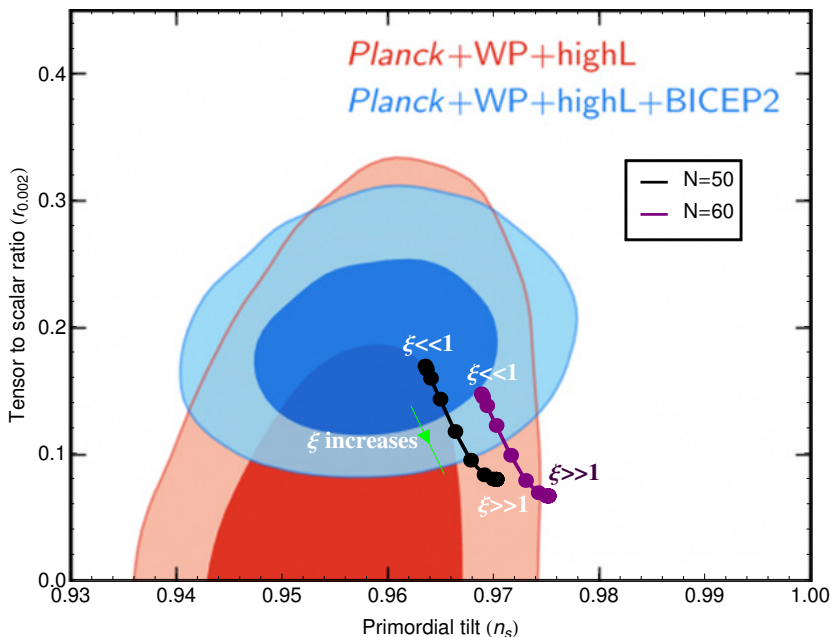
and use eqs. (3.12), (3.16) and (3.19) to express  $n_s$ ,  $r$  and  $|\zeta|^2$  in terms of  $\mathcal{N}$  for  $\xi \ll 1$  case as

$$n_s \simeq 1 - \frac{3}{2\varphi^2} - \frac{3}{16\varphi^2 \ln^2(\varphi/\Lambda)} - \frac{5}{4\varphi^2 \ln(\varphi/\Lambda)} + \mathcal{O}(\xi) \simeq 1 - \frac{3}{\mathcal{N}} + \mathcal{O}(\xi), \quad (4.24)$$

$$r \simeq \frac{8}{\varphi^2} + \frac{1}{2\varphi^2 \ln^2(\varphi/\Lambda)} + \frac{4}{\varphi^2 \ln(\varphi/\Lambda)} + \mathcal{O}(\xi) \simeq \frac{16}{\mathcal{N}} + \mathcal{O}(\xi), \quad (4.25)$$

$$|\zeta|^2 \simeq \frac{8\varphi^6 \ln^3(\varphi/\Lambda)}{3\pi^2(1 + 4 \ln(\varphi/\Lambda))^2} + \mathcal{O}(\xi) \simeq \frac{\mathcal{N}^3}{48\pi^2} + \mathcal{O}(\xi). \quad (4.26)$$

From the above analytical estimations, we see that when  $\xi \gg 1$ ,  $n_s$ ,  $r$  and  $|\zeta|^2$  can satisfy the 95%CL observational bound from Planck data if  $50 < \mathcal{N} < 60$  and  $\xi \sim 10^4$ . Nevertheless,



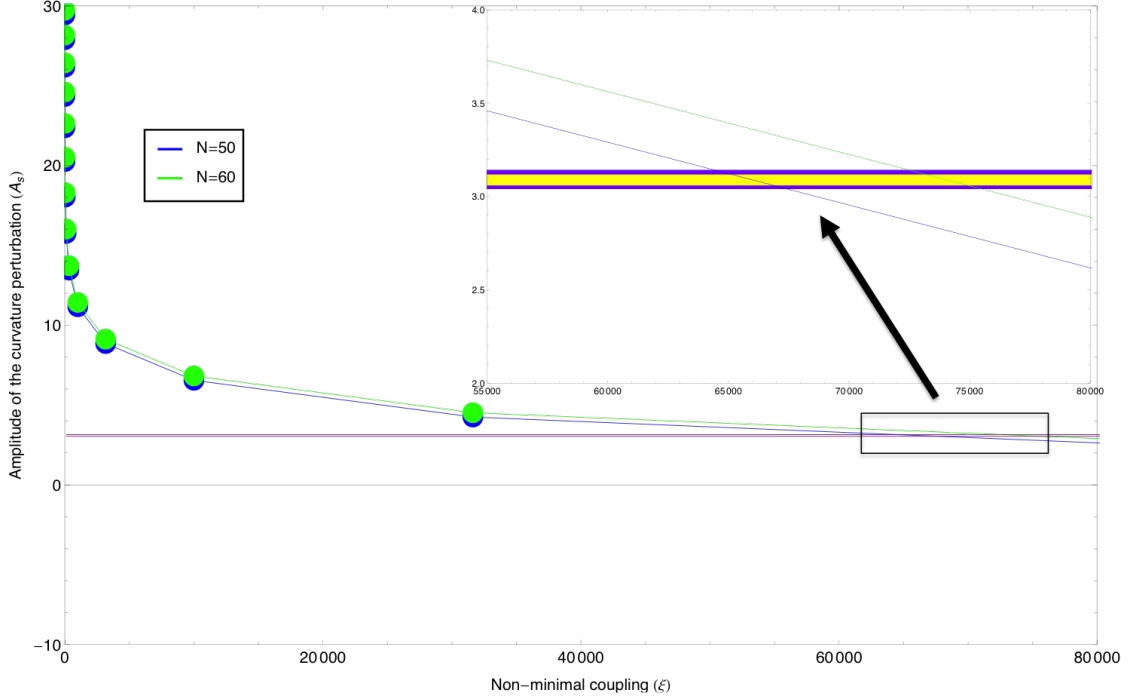
**Figure 3.** The contours display the resulting 68% and 95% confidence regions for the tensor-to-scalar ratio  $r$  and the scalar spectral index  $n_s$ . The red contours are for the Planck+WP+highL data combination, which for this model extension gives a 95% bound  $r < 0.26$  (Planck Collaboration XVI 2013 [21]). The rest represents the BICEP2 constraints on  $r$ . The plots show the numerical data resulting from the GI model by varying the coupling  $\xi$  within a range of  $10^{-3} \lesssim \xi \lesssim 10^6$ .

for such range of  $\mathcal{N}$ ,  $r$  lies outside the  $2\sigma$  CL with BICEP2 results shown in figure 3. The value of  $r$  will increase and then satisfy the bound from BICEP2 results when  $\mathcal{N} \lesssim 45$  or  $\xi \ll 1$ . For the  $\xi \ll 1$  limit, however, it follows from eq. (4.26) that  $|\zeta|^2$  is much larger than the observational bound from Planck data. We see from figure 3 and 4 that these analytical estimations are in good agreement with the values for  $n_s$ ,  $r$  and  $|\zeta|^2$  numerically computed from eqs. (3.11), (3.18) and (3.15). Implementing numerical analysis allows us to explore the behaviour of  $n_s$ ,  $r$  and  $|\zeta|^2$  to what extent they are responsible for the confining scale  $\Lambda$ . From numerical investigations, we come up with the following additional information: firstly, for large  $\xi$  case,  $n_s$  and  $|\zeta|^2$  get higher while  $r$  gets smaller when  $\Lambda$  decreases. Secondly,  $n_s$  and  $r$  get higher when  $\Lambda$  increases towards unity for  $\xi \ll 1$  case.

In general, the potential arises in this model of composite inflation is quite subtle, because it becomes negative when  $\varphi < \Lambda$  and its minimum is also negative. However, the authors of [54] investigated cosmological evolution in models in which the effective potential become negative at some values of the inflaton field. They also discovered several qualitatively new features as compared to those of the positive one. In the next section, we will consider the another compelling paradigm for composite inflation model that leads to a new form of the potential for inflation.

### 4.3 Super-glueball inflation (SgBI)

This model has been explored in [20] in the context of four-dimensional strongly interacting field theories non-minimally coupled to gravity. It is proposed by suitably modifying the theory of an ordinary QCD. We addressed in the appendix C the underlying gauge theory of inflation and showed how the non-minimal coupling comes into play in the theory at the fundamental description. We also constrain such model by using the results from Planck and



**Figure 4.** The plot shows the relation between the amplitude of the power spectrum  $\mathcal{A}_s$  and the non-minimal coupling  $\xi$  within a range of  $10^{-3} \lesssim \xi \lesssim 10^6$  for  $\mathcal{N} = 50, 60$  predicted by the GI model. The horizontal bands represent the  $1\sigma$  (yellow) and  $2\sigma$  (purple) CL for  $\mathcal{A}_s$  obtained from Planck.

BICEP2. The effective Lagrangian for this model has been derived in [20]. Following the literature, we find

$$f(\Phi) = \frac{4\alpha N_c^2}{9} \ln^2(\Phi/\Lambda^3). \quad (4.27)$$

As always investigated in standard fashion, we take the scalar component part of the super-glueball action and coupled it non-minimally to gravity. Focusing only on the modulus of the inflaton field and taking the next step in order to write the non-minimally coupled scalar component part of the super-glueball action to gravity, the resulting action in the Jordan frame reads

$$\mathcal{S}_{\text{SgbI}} = \int d^4x \sqrt{-g} \left[ \frac{1 + N_c^2 \xi \Phi^{2/3}}{2} R - \frac{N_c^2}{2\alpha} \Phi^{-4/3} g^{\mu\nu} \partial_\mu \Phi \partial_\nu \Phi - V_{\text{SgbI}}(\Phi) \right], \quad (4.28)$$

where

$$V_{\text{SgbI}}(\Phi) = \frac{4\alpha N_c^2}{9} \Phi^{4/3} \ln^2(\Phi/\Lambda^3). \quad (4.29)$$

In order to re-write the action in terms of the new field  $\varphi$  possessing unity canonical dimension, we replace the field  $\Phi$  with  $\Phi = \varphi^3$  since in this case  $D = 3$ . Therefore, the action of the theory for our investigation is given by

$$\mathcal{S}_{\text{SgbI}} = \int d^4x \sqrt{-g} \left[ \frac{1 + N_c^2 \xi \varphi^2}{2} R - \frac{9N_c^2}{2\alpha} g^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi - V_{\text{SgbI}}(\varphi) \right], \quad (4.30)$$

where

$$V_{\text{SgbI}}(\varphi) = 4\alpha N_c^2 \varphi^4 (\ln[\varphi/\Lambda])^2, \quad (4.31)$$

with  $N_c$  a number of colours. With the action given above, we find

$$F(\varphi) = 1 + N_c^2 \xi \varphi^2 \quad \text{and} \quad G = \frac{9N_c^2}{\alpha}. \quad (4.32)$$

Using the similar approximations to those of the Glueball Inflation, the number of e-foldings for this inflation model in the large and small  $\xi$  limits are respectively approximated by

$$\mathcal{N} \simeq \frac{3}{2} \left( \ln^2(\varphi/\Lambda) - \ln^2(\varphi_e/\Lambda) \right) + \mathcal{O}(1/\xi), \quad (4.33)$$

and

$$\mathcal{N} \simeq \int_{\varphi_e}^{\varphi} \frac{9N_c^2 \varphi \ln(\varphi/\Lambda) d\varphi}{4\alpha \ln(\varphi/\Lambda) + 2\alpha} + \mathcal{O}(\xi) \simeq \frac{9N_c^2}{8\alpha} (\varphi^2 - \varphi_e^2) + \mathcal{O}(\xi). \quad (4.34)$$

Regarding to the above relations between the number of e-foldings and  $\varphi$ , we can compute  $n_s$ ,  $r$  and  $|\zeta|^2$  in terms of  $\mathcal{N}$  for  $\xi \gg 1$  and  $\xi \ll 1$  limits. For a large  $\xi$  limit, eqs. (3.13), (3.20) and (3.17) yield

$$n_s \simeq 1 - \frac{4}{3 \ln^2(\varphi/\Lambda)} + \mathcal{O}(1/\xi) \simeq 1 - \frac{2}{\mathcal{N}} + \mathcal{O}(1/\xi), \quad (4.35)$$

$$r \simeq \frac{16}{3 \ln^2(\varphi/\Lambda)} + \mathcal{O}(1/\xi) \simeq \frac{8}{\mathcal{N}} + \mathcal{O}(1/\xi), \quad (4.36)$$

$$|\zeta|^2 \simeq \frac{4\alpha \ln^4(\varphi/\Lambda)}{72N_c^2 \pi^2 \xi^2} + \mathcal{O}(1/\xi^3) \simeq \frac{2\alpha N^2}{81N_c^2 \pi^2 \xi^2} + \mathcal{O}(1/\xi^3), \quad (4.37)$$

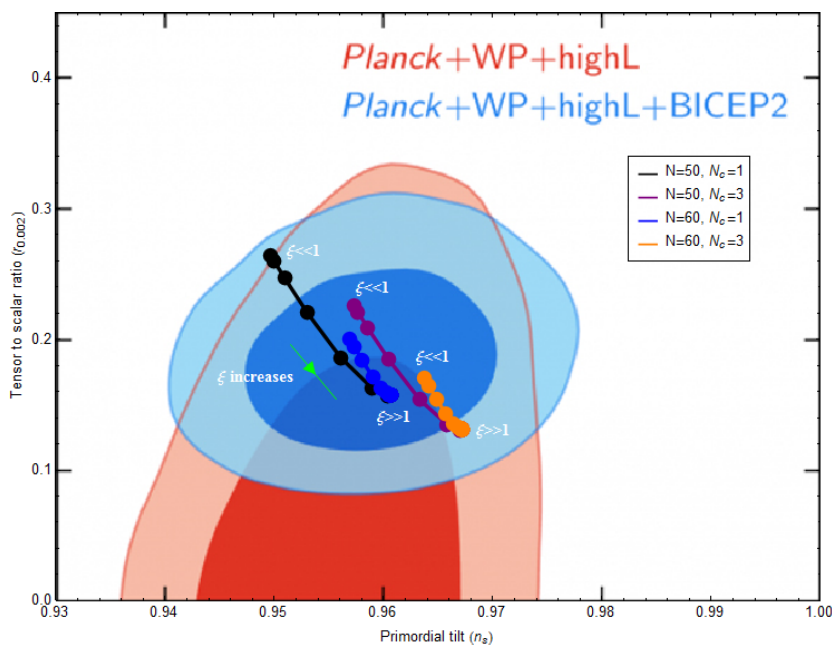
and for a small  $\xi$  limit, eqs. (3.12), (3.19) and (3.16) give

$$n_s \simeq 1 - \frac{24\alpha}{9N_c^2 \varphi^2} - \frac{8\alpha}{9N_c^2 \varphi^2 \ln(\varphi/\Lambda)} - \frac{20\alpha}{9N_c^2 \varphi^2 \ln^2(\varphi/\Lambda)} + \mathcal{O}(\xi) \simeq 1 - \frac{3}{\mathcal{N}} + \mathcal{O}(\xi), \quad (4.38)$$

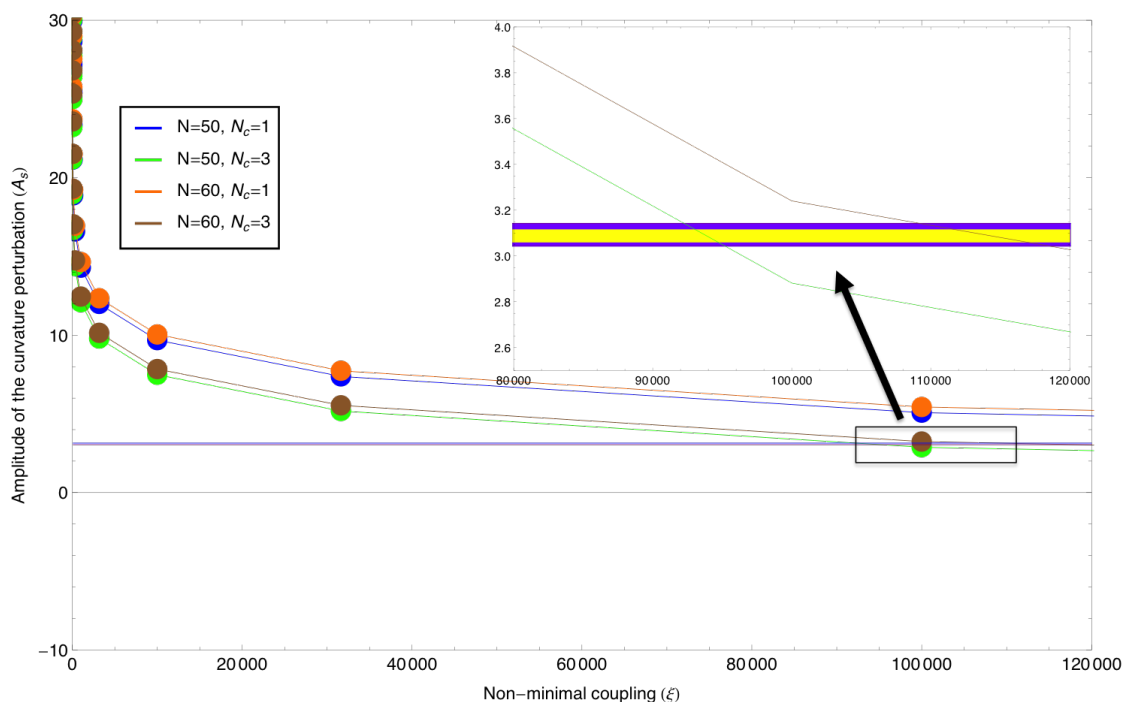
$$r \simeq \frac{32\alpha \left( 1 + 4 \ln(\varphi/\Lambda) + 4 \ln^2(\varphi/\Lambda) \right)}{9N_c^2 \varphi^2 \ln^2(\varphi/\Lambda)} + \mathcal{O}(1/\xi) \simeq \frac{16}{\mathcal{N}} + \mathcal{O}(\xi), \quad (4.39)$$

$$|\zeta|^2 \simeq \frac{N_c^4 \varphi^6 \ln^4(\varphi/\Lambda)}{12\pi^2 \left( 1 + 2 \ln(\varphi/\Lambda) \right)^2} + \mathcal{O}(\xi) \simeq \frac{64N_c \alpha^3 \mathcal{N}^3}{2187\pi^2} + \mathcal{O}(\xi). \quad (4.40)$$

The main results from the above analytical estimations are similar to those of Glueball Inflation. The interesting different result is that for this model of inflation,  $r$  can be large enough to satisfy the bound launched by BICEP2. We compute  $n_s$ ,  $r$  and  $|\zeta|^2$  for this model numerically using eqs. (3.11), (3.18) and (3.15) for a given number of e-foldings, and plot the relevant results in figures 5 and 6. Concretely, the predictions of this model are satisfactorily consistent with the BICEP2 data at 68% CL if  $\mathcal{N} \subseteq [50, 60]$  and  $\xi \sim 10^4 - 10^5$  with an exceptional range of  $\Lambda$ . Regarding to such ranges of  $\mathcal{N}$ ,  $\xi$  and  $\Lambda$ , we discover that the predictions of this model can also satisfy the Planck data. Except for the results illustrated in figures 5 and 6, the numerical implementations also address that, for  $\xi \gg 1$ ,  $r$  gets smaller when  $\Lambda$  decreases, and  $r$  is in tension with the BICEP2 results when  $\Lambda < 10^{-4}$ .

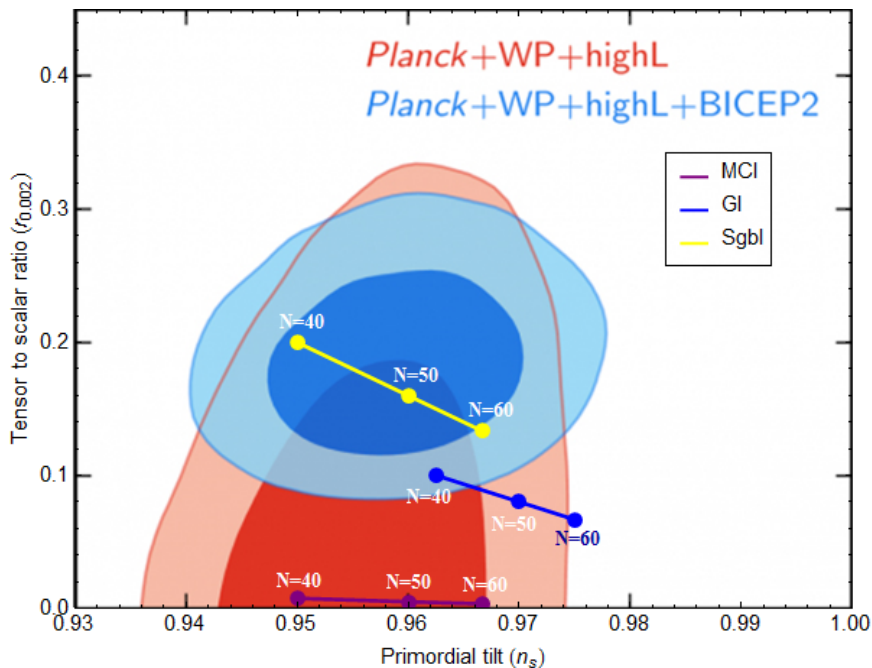


**Figure 5.** The contours display the resulting 68% and 95% confidence regions for the tensor-to-scalar ratio  $r$  and the scalar spectral index  $n_s$ . The red contours are for the Planck+WP+highL data combination, which for this model extension gives a 95% bound  $r < 0.26$  (Planck Collaboration XVI 2013 [21]). The rest represents the BICEP2 constraints on  $r$ . The plots show the numerical data resulting from the SgbI model by varying the coupling  $\xi$ .



**Figure 6.** The plot shows the relation between the amplitude of the power spectrum  $\mathcal{A}_s$  and the non-minimal coupling  $\xi$  for  $\mathcal{N} = 50, 60$  predicted by the SgbI model. The horizontal bands represent the  $1\sigma$  (yellow) and  $2\sigma$  (purple) CL for  $\mathcal{A}_s$  obtained from Planck.





**Figure 7.** The contours display the resulting 68% and 95% confidence regions for the tensor-to-scalar ratio  $r$  and the scalar spectral index  $n_s$ . The red contours are for the Planck+WP+highL data combination, which for this model extension gives a 95% bound  $r < 0.26$  (Planck Collaboration XVI 2013 [21]). The rest represents the BICEP2 constraints on  $r$ . The plots summarise the predictions of composite models examined in this work (MCI, GI, Sgbl) assuming the number of e-foldings  $\mathcal{N}$  to the end of inflation lies in the interval  $[40, 60]$  and  $\xi \gg 1$ .

## 5 Conclusions

In this work, we constrain the model parameters of various composite inflationary models using the observational bound for  $n_s$  and  $r$  from Planck and recent BICEP2 observations, and use  $\mathcal{A}_s$  from Planck data. The general action for the composite inflation has to be in the form of scalar-tensor theory in which the inflaton is non-minimally coupled to gravity. We compute the power spectra for the curvature perturbations by using the usual slow-roll approximations. We summarise our findings, illustrated in figure 7 as follows:

We discover for MCI model with  $\xi \gg 1$  that the predictions lie well inside the joint 68% CL for the Planck+WP+highL data for  $\mathcal{N} = [40, 60]$ , whilst for  $\mathcal{N} = 60$  this model lies on the boundary of  $1\sigma$  region of the Planck+WP+highL data. However, with  $\xi \gg 1$ , the model predictions is in tension with the recent BICEP2 contours. This is so since the model predictions yield quite small values of  $r$ . Concretely, the model predicts  $\epsilon \sim 1/\mathcal{N}^2$  which no longer holds in light of the BICEPS results for  $r = 16\epsilon$  such that  $r = 0.2^{+0.07}_{-0.05}$ . Nevertheless, this tension can be relaxed if  $\xi$  is very small, i.e.  $\xi \sim 10^{-3}$ . If this is the case,  $\mathcal{A}_s$  cannot satisfy the Planck data unless  $\kappa$  gets extremely small, i.e.  $\kappa \sim 10^{-13}$ . Unfortunately, the prediction with very small  $\kappa$  is opposed to the underlying theory. This model predicts  $n_s \simeq 0.960$  and  $r \simeq 0.0048$  for  $\mathcal{N} = 50$  with  $\xi \gg 1$ . Likewise, the Higgs inflation is also in tension with the recent BICEP2 data. The authors in [55] claim the incompatibility between the BICEP2 results and the predictions of Higgs inflation. This tension can be alleviated with the presence of sizable quantum departures from the  $\phi^4$ -inflationary model with the non-minimally coupled scenario [56].

The inflationary observables predicted by the GI model lie well inside the Planck data for  $\mathcal{N} \subseteq [45, 60]$  at the  $2\sigma$  region of the contours for large  $\xi$ . Nevertheless, the predictions of this model with  $\mathcal{N} > 45$  are in tension with the BICEP2 data. However, this can be accommodated to the data allowing that the smaller number of  $\mathcal{N}$  is basically required. It is obvious that  $\mathcal{N}$  is model-dependent quantity. However, it is quite subtle if we have  $\mathcal{N} \lesssim 45$  for model of inflation. This is so since, in order to solve the horizon problem, in the common formulation one frequently use at least  $\mathcal{N} \subseteq [50, 60]$ . We anticipate this can be further verified by studying the reheating effect.

We discover that  $\mathcal{A}_s$  is well consistent with the Planck data up to  $2\sigma$  CL for  $\mathcal{N} = 60$  with  $7.3 \times 10^4 \lesssim \xi \lesssim 7.5 \times 10^4$ , illustrated in figure 4. However,  $\mathcal{A}_s$  does strongly depend on the number of e-foldings implying that the coupling can be lowered (or raised) with changing  $\mathcal{N}$ . This model provides  $n_s \simeq 0.967$  and  $r \simeq 0.089$  for  $\mathcal{N} = 45$  with  $\xi \gg 1$ . In the small  $\xi$  limit, we find that the predictions for  $n_s$  and  $r$  are consistent with Planck data illustrated in figure 3. In the case  $\xi \ll 1$ , we find that the predictions can also satisfy the recent BICEP2 results. However, with a very small  $\xi$  limit, the amplitude of the curvature perturbation  $\mathcal{A}_s$  lie far away from the 95% CL of the Planck data shown in figure 4.

Surprisingly, the SgBI predictions are fully consistent with BICEP2 constraints for  $\mathcal{N} \subseteq [50, 60]$ . Moreover, the model can also be consistent with the Planck contours at  $1\sigma$  CL. We discover that  $\mathcal{A}_s$  is well consistent with the Planck data up to  $2\sigma$  CL for  $\mathcal{N} = 50$  and  $N_c = 3$  with  $9.2 \times 10^4 \lesssim \xi \lesssim 9.5 \times 10^4$ , illustrated in figure 6. This model provides  $n_s \simeq 0.960$  and  $r \simeq 0.16$  for  $\mathcal{N} = 50$  with  $\xi \gg 1$ . However, with the very small  $\xi$ , the amplitude of the curvature perturbation  $\mathcal{A}_s$  cannot get close to the 95% CL of the Planck data shown in figure 4. It would be nice to use the BICEP2 results to constrain  $\Lambda_{\text{SgBI}}$  since the data provides us the lower bound on  $r$ . According to the recent BICEP2 data, we roughly opt  $r \simeq 0.12$  and use  $N_c = 1(3)$  predicting  $\Lambda_{\text{SgBI}} > 10^{-3}(10^{-4})$  which corresponds to, at least, the GUT energy scale in this investigation, in order to satisfy the BICEP2 data at  $1\sigma$  CL. We hope that the future observations will provide significant confirmation for this model.

Another crucial consequence for the model of inflation is the (p)reheating mechanism. We anticipate to investigate this mechanism by following closely references [51–53].

## Acknowledgments

K.K. is supported by Naresuan University through grant R2555C101. P.C. is financially supported by the Institute for the Promotion of Teaching Science and Technology (IPST) under the project of the ‘‘Research Fund for DPST Graduate with First Placement’’ with Grant No.033/2557.

## A Underlying gauge theory of minimal composite inflation

The underlying gauge theory for the technicolor-inspired inflation is the  $SU(N)$  gauge group with  $N_f = 2$  Dirac massless fermions. The two technifermions transform according to the adjoint representation of  $SU(2)$  technicolor (TC) gauge group, called  $SU(2)_{\text{TC}}$ . In order to examine the symmetry properties of the theory, we arrange them by using the Weyl basis in the following vector transformation according to the fundamental representation of  $SU(4)$

group. The field contents in the Weyl basis are

$$\mathcal{Q}^a = \begin{pmatrix} U_L^a \\ D_L^a \\ -i\sigma^2 U_R^{*a} \\ -i\sigma^2 D_R^{*a} \end{pmatrix}, \quad (\text{A.1})$$

where  $U_L$  and  $D_L$  are the left-handed techniup and technidown respectively, and  $U_R$  and  $D_R$  are the corresponding right-handed particles and the upper index  $a = 1, 2, 3$  is the TC index indicating the three dimensional adjoint representation. With the standard breaking to the maximal diagonal subgroup, the  $SU(4)$  global symmetry spontaneously breaks to  $SO(4)$ . Such a breaking is driven by the formation of the following condensate:

$$\langle \mathcal{Q}_i^\alpha \mathcal{Q}_j^\beta \epsilon_{\alpha\beta} \mathcal{E}^{ij} \rangle = -2 \langle \bar{U}_R U_L + \bar{D}_R D_L \rangle, \quad (\text{A.2})$$

where  $i, j = 1, \dots, 4$  denote the components of the tetraplet of  $\mathcal{Q}$ , and  $\alpha, \beta$  indicate the ordinary spin. The  $4 \times 4$  matrix  $\mathcal{E}^{ij}$  is defined in terms of the 2-dimensional identical matrix,  $\mathbb{1}$ , as

$$\mathcal{E} = \begin{pmatrix} 0 & \mathbb{1} \\ \mathbb{1} & 0 \end{pmatrix}, \quad (\text{A.3})$$

with, for example,  $\epsilon_{\alpha\beta} = -i\sigma_{\alpha\beta}^2$  and  $\langle U_L^\alpha U_R^{*\beta} \epsilon_{\alpha\beta} \rangle = -\langle \bar{U}_R U_L \rangle$ . The connection between the composite scalar fields and the underlying technifermions can be obtained from the transformation properties of  $SU(4)$ . To this end, we observe that the elements of the matrix  $\mathcal{M}$  transform like technifermion bilinears such that

$$\mathcal{M}_{ij} \sim \mathcal{Q}_i^\alpha \mathcal{Q}_j^\beta \epsilon_{\alpha\beta} \quad \text{with } i, j = 1, \dots, 4. \quad (\text{A.4})$$

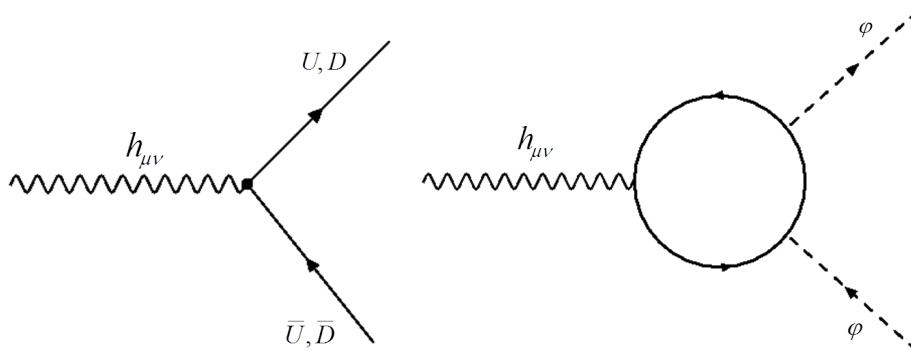
The composite action can be built in terms of the matrix  $\mathcal{M}$  in the Jordan frame as

$$\mathcal{S}_{\text{MCLJ}} = \int d^4x \sqrt{-g} \left[ \frac{M_{\text{P}}^2}{2} R + \frac{1}{2} \xi \text{Tr} [\mathcal{M} \mathcal{M}^\dagger] R + \mathcal{L}_{\text{MWT}} \right], \quad (\text{A.5})$$

where  $\mathcal{L}_{\text{MWT}}$  is the Lagrangian density of the MWT sector, see [18] for more details. The details of this sector are not relevant for the present discussion. From the above action, the non-minimally coupled term corresponds at the fundamental level to a four-fermion interaction term coupled to the Ricci scalar in the following way:

$$\frac{1}{2} \xi \text{Tr} [\mathcal{M} \mathcal{M}^\dagger] R = \frac{1}{2} \xi \frac{(\mathcal{Q}\mathcal{Q})^\dagger \mathcal{Q}\mathcal{Q}}{\Lambda_{\text{Ex}}^4} R, \quad (\text{A.6})$$

where  $\Lambda_{\text{Ex}}$  is a new high energy scale in which this operator generates. Here the non-minimal coupling is added at the fundamental level showing that the non-minimal coupling is well motivated at the level of the fundamental description. Likewise, it seems natural that the conformal transformation used in the transition from the Jordan frame to the Einstein frame should be performed at the level of the fundamental theory. This can be verified by considering the conformal factor like  $F(\mathcal{Q}) \sim (M_{\text{P}}^2 + \xi(\mathcal{Q}\mathcal{Q})^\dagger \mathcal{Q}\mathcal{Q})/M_{\text{P}}^2$ . This leads to the same form of that at the level of the effective low energy description. Moreover, at the level of the fundamental theory, the transition from the Jordan frame to the Einstein frame can be also implemented by redefining the fundamental fields investigated in [52].



**Figure 8.** In the theory with non-minimal coupling to gravity, possible processes show the graviton decays to two technifermions (left) and then such fermions condensate into the composite field (right).

The possible tree-level scattering process features the interaction between graviton and technifermions in a theory with a non-minimal coupling to gravity can be illustrated in figure 8. However, an instructive analysis of the generated coupling of a composite scalar field to gravity has been initiated in the Nambu-Jona-Lasinio (NJL) model [17]. With this regard, the non-minimal coupling apparently seems rather natural. Using the renormalization group equation for the chiral condensate, we find

$$\langle \mathcal{Q}\mathcal{Q} \rangle_{\Lambda_{\text{Ex.}}} \sim \left( \frac{\Lambda_{\text{Ex.}}}{\Lambda_{\text{MCI}}} \right)^\gamma \langle \mathcal{Q}\mathcal{Q} \rangle_{\Lambda_{\text{MCI}}}, \quad (\text{A.7})$$

where the subscripts indicate the energy scale at which the corresponding operators are evaluated, and basically  $\Lambda_{\text{Ex.}} \gg \Lambda_{\text{MCI}}$ . If we assume the fixed value of  $\gamma$  is around two the explicit dependence on the higher energy  $\Lambda_{\text{Ex.}}$  disappears. This is since we have  $\mathcal{M} \sim \langle \mathcal{Q}\mathcal{Q} \rangle_{\Lambda_{\text{MCI}}} / \Lambda_{\text{MCI}}^2$ . According to this model at the effective description, the relevant effective theory consisting of a composite inflaton ( $\varphi$ ) and its pseudo scalar partner ( $\Theta$ ), as well as nine pseudo scalar Goldstone bosons ( $\Pi^{\mathbb{A}}$ ) and their scalar partners ( $\tilde{\Pi}^{\mathbb{A}}$ ) can be assembled in the matrix form such that

$$\mathcal{M} = \left[ \frac{\varphi + i\Theta}{2} + \sqrt{2} \left( i\Pi^{\mathbb{A}} + \tilde{\Pi}^{\mathbb{A}} \right) X^{\mathbb{A}} \right] \mathcal{E}, \quad (\text{A.8})$$

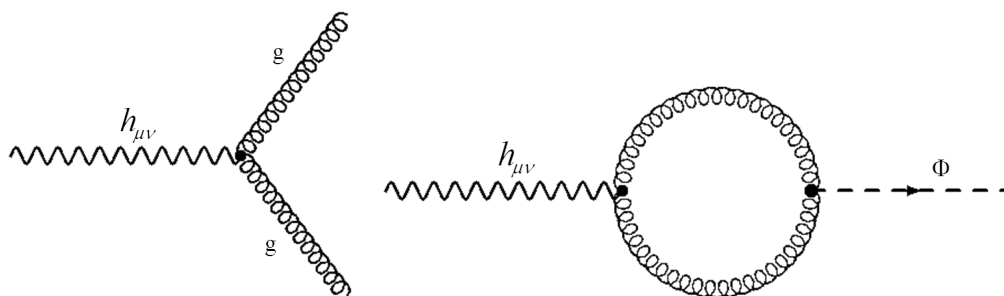
where  $X^{\mathbb{A}}$ 's,  $\mathbb{A} = 1, \dots, 9$ , are the generators of the SU(4) gauge group which do not leave the vacuum expectation value (vev) of  $\mathcal{M}$  invariant, i.e.  $\langle \mathcal{M} \rangle = v\mathcal{E}/2$ ,  $v \equiv \langle \varphi \rangle$ .

## B Underlying gauge theory of glueball inflation

The underlying gauge theory for glueball inflation is the pure SU(N) Yang-Mills gauge theory. The inflaton in this case is the interpolating field describing the lightest glueball. In the same manner with the preceding section, the connection between the composite field and the underlying fundamental description can be also obtained. In this case, the inflaton field is

$$\Phi = \frac{\beta}{\mathfrak{g}} \text{Tr} [\mathcal{G}^{\mu\nu} \mathcal{G}_{\mu\nu}], \quad (\text{B.1})$$

where  $\mathcal{G}_{\mu\nu}$  is the standard non-Abelian field strength,  $\beta$  is the full beta function of the theory in any renormalisation scheme, and  $\mathfrak{g}$  is the gauge coupling. We can also demonstrate that



**Figure 9.** In the theory with non-minimal coupling to gravity, possible processes show the graviton decays to two gluons (left) and then such pure gluons condensate into a glueball (right).

the fundamental degrees of freedom are naturally non-minimally coupled to gravity, and features the description at the fundamental level. In doing so, we introduce the non-minimal coupling term as follows:

$$\xi \left( \frac{\beta}{\mathfrak{g}} \text{Tr} [\mathcal{G}^{\mu\nu} \mathcal{G}_{\mu\nu}] \right)^{1/2} R \equiv \xi \Phi^{1/2} R, \quad (\text{B.2})$$

where the  $\xi$  is a dimensionless quantity. The possible tree-level scattering process features the interaction between graviton and gauge fields in a theory with a non-minimal coupling to gravity can be illustrated in figure 9. Here, together with the preceding section, we have explicitly explained how the introduction of the non-minimal coupling is motivated in a natural way with the underlying fundamental descriptions.

### C Underlying gauge theory of super-glueball inflation

The underlying gauge theory of this model is initiated in [31] based on the following considerations. Let us consider the pure  $N = 1$  supersymmetric Yang-Mills (SYM) gauge theory proposed by suitably modifying that of the ordinary QCD. The theory we are considering is the  $SU(N_c)$  gauge group featuring a one flavor ( $N_f = 1$ ) gauge group with Weyl fermions in the adjoint representation. The Lagrangian can be written as

$$\mathcal{L}_{\text{SYM}} = -\frac{1}{4} \mathcal{G}_{\mu\nu}^a \mathcal{G}^{a,\mu\nu} + \frac{i}{2} \bar{\lambda}^{a,\alpha} \not{D}_{ab} \lambda_\alpha^b + \dots, \quad (\text{C.1})$$

where  $\alpha$  is an ordinary spin,  $a = 1, \dots, N_c^2 - 1$ ,  $\lambda^a$  is the spinor field and  $\mathcal{G}_{\mu\nu}^a$ ,  $\not{D}_{ab}$  are the usual Yang-Mills strength tensor and a covariant derivative, respectively. The dots in principle represent “gauge fixing, ghost terms and auxiliary fields” of those are not relevant for our current discussion. This theory is supersymmetric of an arbitrary  $N_c$ . If a strongly interacting regime takes place, the spinor fields (gluon fields) do condensate into a composite field, called super-glueball, which will be identifies as the inflaton  $\Phi$  in this case. The precise form of the inflaton field is prior given in [20] such that  $\Phi = -3\lambda^{a,\alpha} \lambda_\alpha^a / 64\pi^2 N_c$ . As examined in the previous two examples, the fundamental degrees of freedom are naturally non-minimally coupled to gravity, and features the description at the fundamental level. We start with the introduction of the non-minimal coupling term as follows:

$$\frac{N_c^2}{2} \xi \left( \frac{-3\lambda^{a,\alpha} \lambda_\alpha^a}{64\pi^2 N_c} \right)^{2/3} R \equiv \frac{N_c^2 \xi \Phi^{2/3} R}{2}. \quad (\text{C.2})$$

Again, the  $\xi$  is the dimensionless coupling. We have just demonstrated how the introduction of the non-minimal coupling is motivated in a natural way with the underlying fundamental descriptions.

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