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Observable Gravitational Waves From Kinetically Modified Non-Minimal Inflation

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We consider Supersymmetric (SUSY) and non-SUSY models of chaotic inflation based on the simplest power-law potential with exponents $n = 2$ and 4. We propose a convenient non-minimal coupling to gravity and a non-minimal kinetic term which ensure, mainly for $n = 4$, inflationary observables favored by the BICEP2/*Keck Array* and *Planck* results. Inflation can be attained for subplanckian inflaton values with the corresponding effective theories retaining the perturbative unitarity up to the Planck scale.

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1. Introduction

Kinetically modified *Non-minimal* (chaotic) *inflation* (nMI) [1] is a variant of nMI which arises in the presence of a non-canonical kinetic term for the inflaton ϕ – apart from the non-minimal coupling $f_R(\phi)$ between ϕ and the Ricci scalar curvature, R which is required by definition in any model of nMI [2]. In this talk we focus on inflationary models based on a synergy between f_R and the inflaton potential $V_{\rm CI}$, which are selected [1,3,4] as follows

$$V_{\rm CI}(\phi)=\lambda^2\phi^n/2^{n/2} \quad \text{and} \quad f_R=1+c_R\phi^{n/2} \quad \text{with} \quad n=2,4. \tag{1.1}$$

Below, we first (in Sec. 1.1) briefly review the basic ingredients of nMI in a non-*Supersymmetric* (SUSY) framework and constrain the parameters of the models in Sec. 1.3 taking into account a number of observational and theoretical requirements described in Sec. 1.2. Then (in Sec. 1.4) we focus on the problem with perturbative unitarity that plagues [5,6] these models at the strong coupling and motivate the form of $f_{\rm K}$ analyzed in our work.

Throughout the text, the subscript $,\chi$ denotes derivation *with respect to* (w.r.t) the field χ, charge conjugation is denoted by a star (*) and we use units where the reduced Planck scale $m_{\rm P}=2.43\cdot 10^{18}$ GeV is set equal to unity.

1.1 Coupling non-Minimally the Inflaton to Gravity

The action of the inflaton ϕ in the *Jordan frame* (JF), takes the form:

$$\mathsf{S}=\int d^4x\sqrt{-\mathfrak{g}}\left(-\frac{f_R}{2}R+\frac{f_{\rm K}}{2}g^{\mu\nu}\partial_\mu\phi\partial_\nu\phi-V_{\rm CI}(\phi)\right). \tag{1.2}$$

where $\mathfrak{g}$ is the determinant of the background Friedmann-Robertson-Walker metric, $g^{\mu\nu}$ with signature $(+,-,-,-)$, $\langle f_R\rangle\simeq 1$ to guarantee the ordinary Einstein gravity at low energy and we allow for a kinetic mixing through the function $f_{\rm K}(\phi)$. By performing a conformal transformation [3] according to which we define the *Einstein frame* (EF) metric $\widehat{g}_{\mu\nu}=f_R\,g_{\mu\nu}$ we can write S in the EF as follows

$$\mathsf{S}=\int d^4x\sqrt{-\widehat{\mathfrak{g}}}\left(-\frac{1}{2}\widehat{R}+\frac{1}{2}\widehat{g}^{\mu\nu}\partial_\mu\widehat{\phi}\partial_\nu\widehat{\phi}-\widehat{V}_{\rm CI}(\widehat{\phi})\right), \tag{1.3a}$$

where hat is used to denote quantities defined in the EF. We also introduce the EF canonically normalized field, $\widehat{\phi}$, and potential, $\widehat{V}_{\rm CI}$, defined as follows:

$$\frac{d\widehat{\phi}}{d\phi}=J=\sqrt{\frac{f_{\rm K}}{f_R}+\frac{3}{2}\left(\frac{f_{R,\phi}}{f_R}\right)^2} \quad \text{and} \quad \widehat{V}_{\rm CI}=\frac{V_{\rm CI}}{f_R^2}, \tag{1.3b}$$

where the symbol $,\phi$ as subscript denotes derivation w.r.t the field ϕ. Plugging Eq. (1.1) into Eq. (1.3b), we obtain

$$J^2=\frac{f_{\rm K}}{f_R}+\frac{3n^2c_R^2\phi^{n-2}}{8f_R^2} \quad \text{and} \quad \widehat{V}_{\rm CI}=\frac{\lambda^2\phi^n}{2^{n/2}f_R^2}. \tag{1.4}$$

In the pure nMI [2–4] we take $f_{\rm K}=1$ and, for $c_R\gg 1$, we infer from Eq. (1.3b), that f_R determines the relation between $\widehat{\phi}$ and ϕ and controls the shape of $\widehat{V}_{\rm CI}$ affecting thereby the observational predictions – see below.

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1.2 Inflationary Observables – Constraints

A model of nMI can be qualified as successful, if it can become compatible with the following observational and theoretical requirements:

(i) The number of e-foldings $\widehat{N}_\star$ that the scale $k_\star = 0.05/\mathrm{Mpc}$ experiences during this nMI must to be enough for the resolution of the horizon and flatness problems of standard Big Bang, i.e., [7]

$$\widehat{N}_\star = \int_{\widehat{\phi}_{\rm f}}^{\widehat{\phi}_\star} d\widehat{\phi}\, \frac{\widehat{V}_{\rm CI}}{\widehat{V}_{{\rm CI},\widehat{\phi}}} \simeq 55, \tag{1.5}$$

where $\phi_\star$ $[\widehat{\phi}_\star]$ are the value of ϕ $[\widehat{\phi}]$ when $k_\star$ crosses the inflationary horizon. Also $\phi_{\rm f}$ $[\widehat{\phi}_{\rm f}]$ is the value of ϕ $[\widehat{\phi}]$ at the end of nMI, which can be found, in the slow-roll approximation, from the condition

$$\max\{\widehat{\epsilon}(\phi_{\rm f}), |\widehat{\eta}(\phi_{\rm f})|\} = 1, \quad \text{where}$$

$$\widehat{\epsilon} = \frac{1}{2}\left(\frac{\widehat{V}_{{\rm CI},\widehat{\phi}}}{\widehat{V}_{\rm CI}}\right)^2 = \frac{1}{2J^2}\left(\frac{\widehat{V}_{{\rm CI},\phi}}{\widehat{V}_{\rm CI}}\right)^2 \quad \text{and} \quad \widehat{\eta} = \frac{\widehat{V}_{{\rm CI},\widehat{\phi}\widehat{\phi}}}{\widehat{V}_{\rm CI}} = \frac{1}{J^2}\left(\frac{\widehat{V}_{{\rm CI},\phi\phi}}{\widehat{V}_{\rm CI}} - \frac{\widehat{V}_{{\rm CI},\phi}}{\widehat{V}_{\rm CI}}\frac{J_{,\phi}}{J}\right). \tag{1.6}$$

It is evident from the formulas above that non trivial modifications on $f_{\rm K}$ and thus to J may have an pronounced impact on the parameters above modifying thereby the inflationary observables too.

(ii) The amplitude $A_{\rm s}$ of the power spectrum of the curvature perturbation generated by ϕ at $k_\star$ has to be consistent with data [7], i.e.,

$$\sqrt{A_{\rm s}} = \frac{1}{2\sqrt{3}\,\pi}\,\frac{\widehat{V}_{\rm CI}(\widehat{\phi}_\star)^{3/2}}{|\widehat{V}_{{\rm CI},\widehat{\phi}}(\widehat{\phi}_\star)|} = \frac{1}{2\pi}\sqrt{\frac{\widehat{V}_{\rm CI}(\phi_\star)}{6\widehat{\epsilon}_\star}} \simeq 4.627\cdot 10^{-5}, \tag{1.7}$$

where the variables with subscript $\star$ are evaluated at $\phi = \phi_\star$.

(iii) The remaining inflationary observables (the spectral index $n_{\rm s}$, its running $a_{\rm s}$, and the tensor-to-scalar ratio r) – estimated through the relations:

$$\text{(a) } n_{\rm s} = 1 - 6\widehat{\epsilon}_\star + 2\widehat{\eta}_\star, \quad \text{(b) } a_{\rm s} = \frac{2}{3}\left(4\widehat{\eta}^2 - (n_{\rm s}-1)^2\right) - 2\widehat{\xi}_\star \quad \text{and} \quad \text{(c) } r = 16\widehat{\epsilon}_\star, \tag{1.8}$$

with $\widehat{\xi} = \widehat{V}_{{\rm CI},\widehat{\phi}}\widehat{V}_{{\rm CI},\widehat{\phi}\widehat{\phi}\widehat{\phi}}/\widehat{V}^2_{\rm CI}$ – have to be consistent with the data [7], i.e.,

$$\text{(a) } n_{\rm s} = 0.968 \pm 0.009 \quad \text{and} \quad \text{(b) } r \leq 0.12, \tag{1.9}$$

at 95% *confidence level* (c.l.) – pertaining to the ΛCDM$+r$ framework with $|a_{\rm s}| \ll 0.01$. Although compatible with Eq. (1.9b) the present combined *Planck* and BICEP2/*Keck Array* results [8] seem to favor r's of order 0.01 since $r = 0.048^{+0.035}_{-0.032}$ at 68% c.l. has been reported.

(iv) The effective theory describing nMI has to remains valid up to a UV cutoff scale $\Lambda_{\rm UV}$ to ensure the stability of our inflationary solutions, i.e.,

$$\text{(a) } \widehat{V}_{\rm CI}(\phi_\star)^{1/4} \leq \Lambda_{\rm UV} \quad \text{and} \quad \text{(b) } \phi_\star \leq \Lambda_{\rm UV}. \tag{1.10}$$

It is expected that $\Lambda_{\rm UV} \simeq m_{\rm P}$, contrary to the pure nMI with $c_R \gg 1$ where $\Lambda_{\rm UV} \ll m_{\rm P}$ – see Sec. 1.4.

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1.3 The Two Regimes of Synergistic nMI

The models of nMI based on Eq. (1.1) exhibit the following two regimes:

(i) The *weak* c_R regime with $c_R \ll 1$. In this case from Eq. (1.3b) we find $J \simeq 1/f_R$ and applying Eqs. (1.5) and (1.6), the slow-roll parameters and $\widehat{N}_\star$ read

$$\widehat{\varepsilon} \simeq \frac{n^2}{2\phi^2 f_R}, \;\; \widehat{\eta} \simeq 2\left(1-\frac{1}{n}\right)\widehat{\varepsilon} - \frac{4+n}{2n} c_R \phi^{\frac{n}{2}} \widehat{\varepsilon} \quad \text{and} \quad \widehat{N}_\star \simeq \frac{\phi_\star^2}{2n}. \tag{1.11}$$

Imposing the condition of Eq. (1.6) and solving then the latter equation w.r.t ϕ_* we arrive at

$$\phi_{\rm f} \simeq n/\sqrt{2} \quad \text{and} \quad \phi_\star \simeq \sqrt{2n\widehat{N}_\star}\,. \tag{1.12}$$

Inflation is attained, thus, only for $\phi > 1$. On the other hand, Eq. (1.7) implies

$$\lambda = \sqrt{6A_{\rm s} f_{n\star}}\pi n^{(2-n)/4}/\widehat{N}_\star^{(2+n)/4}\,, \tag{1.13}$$

where $f_{n\star} = f_R(\phi_\star) = 1 + c_R(2n\widehat{N}_\star)^{n/4}$. Applying Eq. (1.8) we find that the inflationary observables are c_R-dependent and can be marginally consistent with Eq. (1.9) – see Sec. 3.2. Indeed,

$$n_{\rm s} = 1 - (4+n+n/f_{n\star})/4\widehat{N}_\star, \;\; r = 4n/f_{n\star}\widehat{N}_\star, \tag{1.14a}$$

$$a_{\rm s} = \left(n^2 - n(n+4)f_{n\star} - 4(n+4)f_{n\star}^2\right)/16 f_{n\star}^2 \widehat{N}_\star^2\,. \tag{1.14b}$$

In the limit $c_R \to 0$ or $f_{n\star} \to 1$ the results of the simplest power-law chaotic inflation – with $f_R = f_{\rm K} = 1$ and $V_{\rm CI}$ given in Eq. (1.1) – are recovered. These are by now disfavored by Eq. (1.9).

(ii) The *strong* c_R regime with $c_R \gg 1$. In this case, from Eq. (1.3b) we find

$$J \simeq \sqrt{3} n c_R \phi^{n/2-1}/2\sqrt{2} f_R \quad \text{and} \quad \widehat{V}_{\rm CI} \simeq \lambda^2/2^{n/2} c_R^2. \tag{1.15}$$

We observe that $\widehat{V}_{\rm CI}$ exhibits an almost flat plateau. From Eqs. (1.5) and (1.6) we find

$$\widehat{\varepsilon} \simeq 4/3c_R^2\phi^n, \;\; \widehat{\eta} \simeq -4/3c_R\phi^{n/2} \quad \text{and} \quad \widehat{N}_\star \simeq 3c_R\phi_*^{n/2}/4\,. \tag{1.16}$$

Therefore, $\phi_{\rm f}$ and $\phi_\star$ are found from the condition of Eq. (1.6) and the last equality above, as follows

$$\phi_{\rm f} = \max\{(4/3c_R^2)^{1/n}, (4/3c_R)^{2/n}\} \quad \text{and} \quad \phi_\star = (4\widehat{N}_\star/3c_R)^{2/n}. \tag{1.17}$$

Consequently, nMI can be achieved even with subplanckian ϕ values for $c_R \gtrsim (4\widehat{N}_*/3)^{2/n}$. Also the normalization of Eq. (1.7) implies the following relation between c_R and λ

$$A_{\rm s}^{1/2} \simeq 2^{-(10+n)/4} \frac{\lambda c_R \phi^n}{\pi f_R}\bigg|_{\phi=\phi_\star} \Rightarrow \lambda \simeq \frac{3\cdot 2^{n/4}}{\widehat{N}_\star}\sqrt{2A_{\rm s}}\,\pi\, c_R\,. \tag{1.18}$$

From Eq. (1.8) we obtain the c_R-independent values for the observables:

$$n_{\rm s} \simeq 1 - 2/\widehat{N}_\star \simeq 0.965, \;\; a_{\rm s} \simeq -2/\widehat{N}_\star^2 \simeq -6.4\cdot 10^{-4} \quad \text{and} \quad r \simeq 12/\widehat{N}_\star^2 \simeq 4\cdot 10^{-3}, \tag{1.19}$$

which are in agreement with Eq. (1.9), although with low enough r values.

1.4 The Ultraviolet (UV) Cut-off Scale

In the highly predictive regime with large $c_{\rm K}$, the models violate perturbative unitarity for $n>2$. To see this we analyze the small-field behavior of the theory in order to extract the UV cut-off scale $\Lambda_{\rm UV}$. The result depends crucially on the value of J in Eq. (1.3b) in the vacuum, $\langle\phi\rangle=0$. Namely we have

$$\langle J\rangle=\begin{cases}\sqrt{3/2}c_R & \text{for } n=2,\\ 1 & \text{for } n\neq 2\,.\end{cases} \tag{1.20}$$

For $n=2$ and any c_R we obtain $\widehat\phi\neq\phi$. Expanding the second and third term of S in the right-hand side of Eq. (1.3a) about $\langle\phi\rangle=0$ in terms of $\widehat\phi$ we obtain:

$$J^2\dot\phi^2=\left(1-\sqrt{\frac{8}{3}}\widehat\phi+2\widehat\phi^2-\cdots\right)\dot{\widehat\phi}^2\quad\text{and}\quad \widehat V_{\rm CI}=\frac{\lambda^2\widehat\phi^2}{3c_R^2}\left(1-\sqrt{\frac{8}{3}}\widehat\phi+2\widehat\phi^2-\cdots\right). \tag{1.21}$$

As a consequence $\Lambda_{\rm UV}=m_{\rm P}$ since the expansions above are c_R independent. On the contrary, for $n>2$ we have $\widehat\phi=\phi$ and the expansions of the same terms in Eq. (1.3a) are c_R dependent:

$$J^2\dot\phi^2=\left(1-c_R\widehat\phi^{\frac{n}{2}}+\frac{3n^2}{8}c_R^2\widehat\phi^{n-2}+c_R^2\widehat\phi^{n}-\cdots\right)\dot{\widehat\phi}^2; \tag{1.22a}$$

$$\widehat V_{\rm CI}=\frac{\lambda^2\widehat\phi^n}{2}\left(1-2c_R\widehat\phi^{\frac{n}{2}}+3c_R^2\widehat\phi^n-4c_R^3\widehat\phi^{\frac{3n}{2}}+\cdots\right)\cdot \tag{1.22b}$$

Since the term which yields the smallest denominator for $c_R>1$ is $3n^2c_R^2\widehat\phi^{n-2}/8$ we find [5,6]:

$$\Lambda_{\rm UV}=m_{\rm P}/c_R^{2/(n-2)}\ll m_{\rm P}\,. \tag{1.23}$$

However, if we introduce a *non-canonical kinetic mixing* of the form

$$f_{\rm K}(\phi)=c_{\rm K}f_R^m\quad\text{where}\quad c_{\rm K}=(c_R/r_{R\rm K})^{4/n}\quad\text{and}\quad m\geq 0, \tag{1.24}$$

no problem with the perturbative unitarity emerges for $r_{R\rm K}\leq 1$, even if c_R and/or $c_{\rm K}$ are large – the latter situation is expected if we wish to achieve efficient nMI with $\phi\leq 1$. E.g., for $m=0$ the expansions in Eqs. (1.22a) and (1.22b) can be rewritten replacing c_R with $r_{R\rm K}$ and λ with $\lambda/c_{\rm K}^{n/4}$ – similar expressions can be obtained for other m, too. In other words, the perturbative unitarity can be preserved up to $m_{\rm P}$ if we select a non-trivial $f_{\rm K}$ such that $\langle J\rangle\neq 1$. This requirement lets a functional uncertainty as regards the form of $f_{\rm K}$ during nMI which can be parameterized as shown in Eq. (1.24) given that $\langle f_R\rangle\simeq 1$ – see Sec. 1.1.

We below describe a possible formulation of this type of nMI in the context of *Supergravity* (SUGRA) – see Sec. 2 – and we then analyze the inflationary behavior of these models in Sec. 3. We conclude summarizing our results in Sec. 4.

2. Supergravity Embeddings

The models above – defined by Eqs. (1.1) and (1.24) – can be embedded in SUGRA if we use two gauge singlet chiral superfields $z^\alpha=\Phi,S$, with Φ ($\alpha=1$) and S ($\alpha=2$) being the inflaton and a "stabilizer" field respectively. The EF action for z^α's can be written as [9]

$$\mathsf{S}=\int d^4x\sqrt{-\widehat{\mathfrak{g}}}\left(-\frac{1}{2}\widehat R+K_{\alpha\bar\beta}\widehat g^{\mu\nu}\partial_\mu z^\alpha\partial_\nu z^{*\bar\beta}-\widehat V\right), \tag{2.1a}$$

where summation is taken over the scalar fields z^α, K is the Kähler potential with $K_{\alpha\bar\beta}=K_{,z^\alpha z^{*\bar\beta}}$ and $K^{\alpha\bar\beta}K_{\bar\beta\gamma}=\delta^\alpha_\gamma$. Also $\widehat V$ is the EF F–term SUGRA potential given by

$$\widehat V=e^K\left(K^{\alpha\bar\beta}D_\alpha W D^*_{\bar\beta}W^*-3|W|^2\right), \tag{2.1b}$$

where $D_\alpha W=W_{,z^\alpha}+K_{,z^\alpha}W$ with W being the superpotential. Along the inflationary track determined by the constraints

$$S=\Phi-\Phi^*=0,\ \text{ or }\ s=\bar s=\theta=0 \tag{2.2}$$

if we express Φ and S according to the parametrization

$$\Phi=\phi\, e^{i\theta}/\sqrt2\quad\text{and}\quad S=(s+i\bar s)/\sqrt2, \tag{2.3}$$

$V_{\rm CI}$ in Eq. (1.1) can be produced, in the flat limit, by

$$W=\lambda S\Phi^{n/2}. \tag{2.4}$$

SUPERFIELDS:	S	Φ
$U(1)_R$	1	0
$U(1)$	-1	$2/n$

Table 1: *Charge assignments of the superfields.*

The form of W can be uniquely determined if we impose an R and a global $U(1)$ symmetry with charge assignments shown in Table 1.

On the other hand, the derivation of $\widehat V_{\rm CI}$ in Eq. (1.4) via Eq. (2.1b) requires a judiciously chosen K. Namely, along the track in Eq. (2.2) the only surviving term in Eq. (2.1b) is

$$\widehat V_{\rm CI}=\widehat V(\theta=s=\bar s=0)=e^K K^{SS^*}\,|W_{,S}|^2. \tag{2.5}$$

The incorporation f_R in Eq. (1.1) and $f_{\rm K}$ in Eq. (1.24) dictates the adoption of a logarithmic K [9] including the functions

$$F_R(\Phi)=1+2^{\frac n4}\Phi^{\frac n2}c_R,\ F_{\rm K}=(\Phi-\Phi^*)^2\quad\text{and}\quad F_S=|S|^2-k_S|S|^4. \tag{2.6}$$

Here, F_R is an holomorphic function reducing to f_R, along the path in Eq. (2.2), $F_{\rm K}$ is a real function which assists us to incorporate the non-canonical kinetic mixing generating by $f_{\rm K}$ in Eq. (1.24), and F_S provides a typical kinetic term for S, considering the next-to-minimal term for stability/heaviness reasons [9]. Indeed, $F_{\rm K}$ lets intact $\widehat V_{\rm CI}$, since it vanishes along the trajectory in Eq. (2.2), but it contributes to the normalization of Φ. Taking for consistency all the possible terms up to fourth order, K is written as

$$K_1=-3\ln\left(\frac12(F_R+F_R^*)+\frac{c_{\rm K}}{3\cdot2^{m+1}}(F_R+F_R^*)^mF_{\rm K}-\frac13F_S+\frac{k_\Phi}{6}F_{\rm K}^2-\frac{k_{S\Phi}}{3}F_{\rm K}|S|^2\right). \tag{2.7a}$$

Alternatively, if we do not insist on a pure logarithmic K, we could also adopt the form

$$K_2=-3\ln\left(\frac12(F_R+F_R^*)-\frac13F_S\right)-\frac{c_{\rm K}}{2^m}\frac{F_{\rm K}}{(F_R+F_R^*)^{1-m}}. \tag{2.7b}$$

Moreover, if we place F_S outside the argument of the logarithm similar results are obtained by the following K's – not mentioned in Ref. [1]:

$$K_3=-2\ln\left(\frac12(F_R+F_R^*)+\frac{c_{\rm K}}{2^{m+2}}(F_R+F_R^*)^mF_{\rm K}\right)+F_S, \tag{2.7c}$$

$$K_4=-2\ln\frac{F_R+F_R^*}{2}-\frac{c_{\rm K}}{2^m}\frac{F_{\rm K}}{(F_R+F_R^*)^{1-m}}+F_S. \tag{2.7d}$$

FIELDS	EINGESTATES	MASSES SQUARED			
		SYMBOL	$K=K_1$	$K=K_2$	$K=K_{i+2}$
2 real scalars	$\widehat{\theta}$	$\widehat{m}^2_{\theta}$	$4\widehat{H}^2_{\rm CI}$	$6\widehat{H}^2_{\rm CI}$	
1 complex scalar	$\widehat{s},\widehat{\bar{s}}$	$\widehat{m}^2_{s}$	$6\left(2k_Sf_R-1/3\right)\widehat{H}^2_{\rm CI}$		$12k_S\widehat{H}^2_{\rm CI}$
4 Weyl spinors	$\widehat{\psi}_{\pm}$	$\widehat{m}^2_{\psi\pm}$	$3n^2\widehat{H}^2_{\rm CI}/2c_{\rm K}\phi^2f_R^{1+m}$		

Table 2: Mass-squared spectrum for $K=K_i$ and $K=K_{i+2}$ $(i=1,2)$ along the path in Eq. (2.2).

Note that for $m=0$ [$m=1$], F_R and $F_{\rm K}$ in K_1 and K_3 [K_2 and K_4] are totally decoupled, i.e. no higher order term is needed. Also we use only integer prefactors for the logarithms avoiding thereby any relevant tuning – cf. Ref. [10]. Our models, for $c_{\rm K}\gg c_R$, are completely natural in the 't Hooft sense because, in the limits $c_R\to 0$ and $\lambda\to 0$, the theory enjoys the enhanced symmetries

$$\Phi\to\ \Phi^*,\ \Phi\to\ \Phi+c\quad\text{and}\quad S\to\ e^{i\alpha}S, \tag{2.8}$$

where c is a real number. It is evident that our proposal is realized more attractively within SUGRA than within the non-SUSY set-up, since both $f_{\rm K}$ and f_R originate from the same function K.

To verify the appropriateness of K's in Eqs. (2.7a) – (2.7d), we can first remark that, along the trough in Eq. (2.2), these are diagonal with non-vanishing elements K_{SS^*} and $K_{\Phi\Phi^*}=J^2$, where J is given by Eq. (1.4) for $K=K_i$ and Eq. (1.4) replacing $3/8$ by $1/4$ for $K=K_{i+2}$. Substituting into Eq. (2.5) $K^{SS^*}=1/K_{SS^*}$ and $\exp K=1/f_R^N$, where

$$K_{SS^*}=\begin{cases}1/f_R\\ 1\end{cases}\text{and}\ \ N=\begin{cases}3\\ 2\end{cases}\text{for}\ \ K=\begin{cases}K_i\\ K_{i+2}\end{cases}\text{with}\ \ i=1,2, \tag{2.9}$$

we easily deduce that $\widehat{V}_{\rm CI}$ in Eq. (1.4) is recovered. If we perform the inverse of the conformal transformation described in Eqs. (1.3a) and (1.2) with frame function $\Omega/N=-e^{-K/N}$ we can easily show that $f_R=-\Omega/N$ along the path in Eq. (2.2). Note, finally, that the conventional Einstein gravity is recovered at the SUSY vacuum, $\langle S\rangle=\langle\Phi\rangle=0$, since $\langle f_R\rangle\simeq 1$.

Defining the canonically normalized fields via the relations $d\widehat{\phi}/d\phi=\sqrt{K_{\Phi\Phi^*}}=J$, $\widehat{\theta}=J\theta\phi$ and $(\widehat{s},\widehat{\bar{s}})=\sqrt{K_{SS^*}}(s,\bar{s})$ we can verify that the configuration in Eq. (2.2) is stable w.r.t the excitations of the non-inflaton fields. Taking the limit $c_{\rm K}\gg c_R$ we find the expressions of the masses squared $\widehat{m}^2_{\chi^\alpha}$ (with $\chi^\alpha=\theta$ and s) arranged in Table 2, which approach rather well the quite lengthy, exact formulas. From these expressions we appreciate the role of $k_S>0$ in retaining positive $\widehat{m}^2_s$. Also we confirm that $\widehat{m}^2_{\chi^\alpha}\gg\widehat{H}^2_{\rm CI}=\widehat{V}_{\rm CI0}/3$ for $\phi_{\rm f}\leq\phi\leq\phi_\star$. In Table 2 we display the masses $\widehat{m}^2_{\psi^\pm}$ of the corresponding fermions too with eignestates $\widehat{\psi}_\pm=(\widehat{\psi}_\Phi\pm\widehat{\psi}_S)/\sqrt{2}$, defined in terms of $\widehat{\psi}_S=\sqrt{K_{SS^*}}\psi_S$ and $\widehat{\psi}_\Phi=\sqrt{K_{\Phi\Phi^*}}\psi_\Phi$, where ψ_Φ and ψ_S are the Weyl spinors associated with S and Φ respectively. Note, finally, that $\widehat{m}_{\chi^\alpha}\ll m_{\rm P}$, for any χ^α, contrary to similar cases [11] where the inflaton belongs to gauge non-singlet superfields.

Inserting the derived mass spectrum in the well-known Coleman-Weinberg formula, we can find the one-loop radiative corrections, $\Delta\widehat{V}_{\rm CI}$ to $\widehat{V}_{\rm CI}$. It can be verified that our results are immune from $\Delta\widehat{V}_{\rm CI}$, provided that the renormalization group mass scale Λ, is determined conveniently and $k_{S\Phi}$ and k_S are confined to values of order unity.

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3. Results

The present inflationary scenario depends on the parameters: $n,\ m,\ r_{R\mathrm{K}},\ \lambda/c_{\mathrm{K}}^{n/4}$. Note that the two last combinations of parameters above replace c_{K}, c_R and λ. This is because, if we perform a rescaling $\phi=\tilde\phi/\sqrt{c_{\mathrm{K}}}$, Eq. (1.2) preserves its form replacing ϕ with $\tilde\phi$ and f_{K} with f_R^m where f_R and V_{CI} take, respectively, the forms

$$f_R=1+r_{R\mathrm{K}}\tilde\phi^{n/2}\quad\text{and}\quad V_{\mathrm{CI}}=\lambda^2\tilde\phi^n/2^{n/2}c_{\mathrm{K}}^{n/2},\tag{3.1}$$

which, indeed, depend only on $r_{R\mathrm{K}}$ and $\lambda^2/c_{\mathrm{K}}^{n/2}$. Imposing the restrictions of Sec. 1.2 we can delineate the allowed region of these parameters. Below we first extract some analytic expressions – see Sec. 3.1 – which assist us to interpret the exact numerical results presented in Sec. 3.2.

3.1 Analytic Results

Assuming $c_{\mathrm{K}}\gg c_R$, Eq. (1.3b) yields $J\simeq\sqrt{c_{\mathrm{K}}}/f_R^{(1-m)/2}$. Inserting the last one and $\widehat V_{\mathrm{CI}}$ from Eq. (1.1) in Eq. (1.6) we extract the slow-roll parameters for this model as follows – cf. Eq. (1.11):

$$\widehat\varepsilon=n^2/2\phi^2c_{\mathrm{K}}f_R^{1+m}\quad\text{and}\quad\widehat\eta=2\,(1-1/n)\,\widehat\varepsilon-(4+n(1+m))\,c_R\phi^{n/2}\widehat\varepsilon/2n\,.\tag{3.2}$$

Given that $\phi\ll1$ and so $f_R\simeq1$, nMI terminates for $\phi=\phi_{\mathrm{f}}$ found by the condition

$$\phi_{\mathrm{f}}\simeq\max\{n/\sqrt{2c_{\mathrm{K}}},\sqrt{(n-1)n/c_{\mathrm{K}}}\}\,,\tag{3.3}$$

in accordance with Eq. (1.6). Since $\phi_\star\gg\phi_{\mathrm{f}}$, from Eq. (1.5) we find

$$\widehat N_\star=\frac{c_{\mathrm{K}}\phi_\star^2}{2n}\,{}_2F_1\left(-m,4/n;1+4/n;-c_R\phi_\star^{n/2}\right)=\begin{cases}c_{\mathrm{K}}\phi_\star^2/2n & \text{for } m=0,\\ (f_R^{1+m}-1)/8(1+m)r_{R\mathrm{K}} & \text{for } n=4,\end{cases}\tag{3.4}$$

where ${}_2F_1$ is the Gauss hypergeometric function. Concentrating on the cases with $m=0$ or $n=4$, we solve Eq. (3.4) w.r.t $\phi_\star$ with results

$$\phi_\star\simeq\begin{cases}\sqrt{2n\widehat N_\star/c_{\mathrm{K}}} & \text{for}\quad m=0,\\ \sqrt{f_{m\star}-1}/\sqrt{r_{R\mathrm{K}}c_{\mathrm{K}}} & \text{for}\quad n=4,\end{cases}\tag{3.5}$$

where $f_{m\star}^{1+m}=1+8(m+1)r_{R\mathrm{K}}\widehat N_\star$. In both cases there is a lower bound on c_{K}, above which $\phi_\star<1$ and so, our proposal can be stabilized against corrections from higher order terms – e.g., for $n=4, m=1$ and $r_{R\mathrm{K}}=0.03$ we obtain $140\lesssim c_{\mathrm{K}}\lesssim1.4\cdot10^6$ for $3.3\cdot10^{-4}\lesssim\lambda\lesssim3.5$. The correlation between λ and $c_{\mathrm{K}}^{n/4}$ can be found from Eq. (1.7). For $m=0$ this is given by Eq. (1.13) replacing λ with $\lambda/c_{\mathrm{K}}^{n/4}$ and c_R with $r_{R\mathrm{K}}$ in the definition of $f_{n\star}$. For $n=4$ we obtain

$$\lambda=16\sqrt{3A_{\mathrm{s}}}\,\pi\,c_{\mathrm{K}}\,r_{R\mathrm{K}}^{3/2}/(f_{m\star}-1)^{3/2}f_{m\star}^{(1+m)/2}\,.\tag{3.6}$$

As regards the inflationary observables, these are obviously given by Eqs. (1.14a) and (1.14b) for the trivial case with $m=0$. For $m\neq0$, however, these are heavily altered. In particular, for $n=4$ we obtain

$$n_{\mathrm{s}}=1-8r_{R\mathrm{K}}\frac{m-1-(m+2)f_{m\star}}{(f_{m\star}-1)f_{m\star}^{1+m}},\ \ r=\frac{128r_{R\mathrm{K}}}{(f_{m\star}-1)f_{m\star}^{1+m}},\tag{3.7a}$$

$$a_{\mathrm{s}}=\frac{64r_{R\mathrm{K}}^2(1+m)(m+2)}{(f_{m\star}-1)^2f_{m\star}^{4(1+m)}}f_{m\star}^2\left(f_{m\star}^{2m}\left(\frac{1-m}{m+2}+\frac{2m-1}{m+1}f_{m\star}\right)-f_{m\star}^{2(1+m)}\right).\tag{3.7b}$$

The formulae above is valid only for $r_{R\mathrm{K}}>0$ – see Eq. (3.5) – and is simplified [1] for low m's.

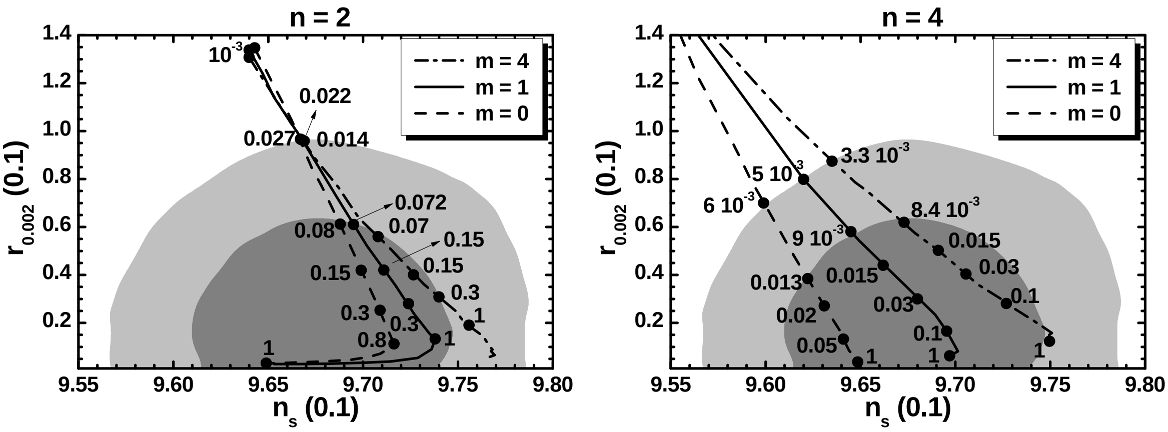

Figure 1: *Allowed curves in the* $n_{\rm s}-r_{0.002}$ *plane for* $n=2$ *and* 4, $m=0$ *(dashed lines),* $m=1$ *(solid lines),* $m=4$ *(dot-dashed lines), and various* $r_{R\rm K}$*'s indicated on the curves. The marginalized joint* 68% *[*95%*] regions from Planck,* BICEP2*/Keck Array and BAO data are depicted by the dark [light] shaded contours.*

3.2 Numerical Results

The conclusions obtained in Sec. 3.1 can be verified and extended to others n's and m's numerically. In particular, enforcing Eqs. (1.5) and (1.7) we can restrict $\phi_\star$ and $\lambda/c_{\rm K}^{n/4}$. Then we can compute the model predictions via Eq. (1.8), for any selected m,n and $r_{R\rm K}$. The outputs, encoded as lines in the $n_{\rm s}-r_{0.002}$ plane, are compared against the observational data [7, 8] in Fig. 1 for $n=2$ (left panel) and 4 (right panel) setting $m=0,1$ and 4 – dashed, solid, and dot-dashed lines respectively. The variation of $r_{R\rm K}$ is shown along each line. To obtain an accurate comparison, we compute $r_{0.002}=16\widehat{\varepsilon}(\phi_{0.002})$ where $\phi_{0.002}$ is the value of ϕ when the scale $k=0.002/{\rm Mpc}$, which undergoes $\widehat{N}_{0.002}=(\widehat{N}_\star+3.22)$ e-foldings during nMI, crosses the inflationary horizon.

From the plots in Fig. 1 we observe that, for low enough $r_{R\rm K}$'s – i.e. $r_{R\rm K}=10^{-4}$ and 0.001 for $n=4$ and 2 –, the various lines converge to the $(n_{\rm s},r_{0.002})$'s obtained within the simplest models of chaotic inflation with the same n. At the other end, the lines for $n=4$ terminate for $r_{R\rm K}=1$, beyond which the theory ceases to be unitarity safe – as anticipated in Sec. 1.4 – whereas the $n=2$ lines approach an attractor value, comparable with the value in Eq. (1.19), for any m.

For $m=0$ we reveal the results of Sec. 1.3, i.e. the displayed lines are almost parallel for $r_{0.002}\geq 0.02$ and converge at the values in Eq. (1.19) – for $n=4$ this is reached even for $r_{R\rm K}=1$. Our estimations in Eqs. (1.14a) – (1.14b) are in agreement with the numerical results for $n=2$ and $r_{R\rm K}\lesssim 1$ or $n=4$ and $r_{R\rm K}\lesssim 0.05$. We observe that the $n=2$ line is closer to the central values in Eq. (1.9) whereas the $n=4$ one deviates from those.

For $m>0$ the curves change slopes w.r.t to those with $m=0$ and move to the right. As a consequence, for $n=4$ they span densely the 1-σ ranges in Eq. (1.9) for quite natural $r_{R\rm K}$'s – e.g. $0.005\lesssim r_{R\rm K}\lesssim 0.1$ for $m=1$. It is worth mentioning that the requirement $r_{R\rm K}\leq 1$ (for $n=4$) provides a lower bound on $r_{0.002}$, which ranges from 0.004 for $m=0$ to 0.015 (for $m=4$). Therefore, our results are testable in the forthcoming experiments [12] hunting for primordial gravitational waves. Note, finally, that our findings in Eqs. (3.7a) – (3.7b) approximate fairly the numerical outputs for $0.003\lesssim r_{R\rm K}\leq 1$.

4. Conclusions

We reviewed the implementation of kinetically modified nMI in both a non-SUSY and a SUSY framework. The models are tied to the potential $V_{\rm CI}$ and the coupling function of the inflaton to gravity given in Eq. (1.1) and the non-canonical kinetic mixing in Eq. (1.24). This setting can be elegantly implemented in SUGRA too, employing the super-and Kähler potentials given in Eqs. (2.4) and (2.7a) – (2.7d). Prominent in this realization is the role of a shift-symmetric quadratic function $F_{\rm K}$ in Eq. (2.6) which remains invisible in the SUGRA scalar potential while dominates the canonical normalization of the inflaton. Using $m \geq 0$ and confining $r_{R\rm K}$ to the range $(3.3 \cdot 10^{-3} - 1)$, where the upper bound does not apply to the $n = 2$ case, we achieved observational predictions which may be tested in the near future and converge towards the "sweet" spot of the present data – especially for $n = 4$. These solutions can be attained even with subplanckian values of the inflaton requiring large $c_{\rm K}$'s and without causing any problem with the perturbative unitarity. It is gratifying, finally, that the most promising case of our proposal with $n = 4$ can be studied analytically and rather accurately.

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References

[1] C. Pallis, *Phys. Rev. D* **91**, 123508 (2015) [arXiv:1503.05887].

[2] D. S. Salopek, J. R. Bond and J.M. Bardeen, *Phys. Rev. D* **40**, 1753 (1989);
F.L. Bezrukov and M. Shaposhnikov, *Phys. Lett. B* **659**, 703 (2008) [arXiv:0710.3755].

[3] C. Pallis, *Phys. Lett. B* **692**, 287 (2010) [arXiv:1002.4765];
C. Pallis and Q. Shafi, *Phys. Rev. D* **86**, 023523 (2012) [arXiv:1204.0252].

[4] R. Kallosh, A. Linde, and D. Roest, *Phys. Rev. Lett.* **112**, 011303 (2014) [arXiv:1310.3950].

[5] J.L.F. Barbon and J.R. Espinosa, *Phys. Rev. D* **79**, 081302 (2009) [arXiv:0903.0355];
C.P. Burgess, H.M. Lee, and M. Trott, *JHEP* **07**, 007 (2010) [arXiv:1002.2730].

[6] A. Kehagias, A.M. Dizgah, and A. Riotto, *Phys. Rev. D* **89**, 043527 (2014) [arXiv:1312.1155].

[7] *Planck* Collaboration, arXiv:1502.02114.

[8] P.A.R. Ade *et al.*, *Phys. Rev. Lett.* **114**, 101301 (2015) [arXiv:1502.00612].

[9] M.B. Einhorn and D.R.T. Jones, *JHEP* **03**, 026 (2010) [arXiv:0912.2718];
H.M. Lee, *JCAP* **08**, 003 (2010) [arXiv:1005.2735];
S. Ferrara *et al.*, *Phys. Rev. D* **83**, 025008 (2011) [arXiv:1008.2942];
C. Pallis and N. Toumbas, *JCAP* **02**, 019 (2011) [arXiv:1101.0325].

[10] C. Pallis, *JCAP* **10**, 058 (2014) [arXiv:1407.8522];
C. Pallis and Q. Shafi, *JCAP* **03**, 023 (2015) [arXiv:1412.3757].

[11] G. Lazarides and C. Pallis, arXiv:1508.06682;
C. Pallis, to appear.

[12] P. Creminelli *et al.*, arXiv:1502.01983.