Higher Derivative D-term Inflation in New-minimal Supergravity

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Abstract

We revisit the D-term inflation and amend it with ghost-free higher derivative couplings of chiral superfields to super-curvature. These couplings realize a more generic inflationary phase in supergravity. After pointing out that a consistent embedding of these specific higher derivatives is known to exist only in the new-minimal supergravity, we show how a potential for the scalar component may arise due to a Fayet-Iliopoulos D-term. We then turn to inflationary cosmology and explicitly discuss different types of potentials, which capture properties of the common scenarios. These models thanks to the derivative coupling: i) naturally evade the supergravity $\eta$-problem, ii) drive inflation for a wider range of parameter values and, iii) can fit the Planck data which, for particular models, would be impossible otherwise.
1 Introduction

Observational data strongly indicate that an inflationary phase did occur at some stage in the early universe. The inflationary scenario provides the most convincing answer to the problems of homogeneity, isotropy and flatness, while it also explains the origin of the large scale structure [1, 2]. The source of inflation can be effectively described by a scalar field, however, its true nature and origin are still unclear. In the case of a fundamental scalar field, theories which support their existence are well-motivated. In fact supersymmetry naturally accommodates scalar particles and offers a deeper theoretical understanding to their existence. The framework under which supersymmetric theories during inflation should be studied is supergravity [3].

In a supergravity theory, which is a non-renormalizable theory, the number of couplings that need to be specified is in principle infinite. The cut-off of the theory is considered to be the Planck mass, $M_P$, and an estimation in supergravity ceases to be reliable for field values $\phi \gg M_P$ where the non-renormalizable terms have to be taken into account. For example in the popular chaotic models, based on potentials $V \propto \phi^n$, super-Planckian values are necessary for sufficient inflation. Even though one can construct models in which the inflaton field value experiences sub-Planckian variations, a generic supergravity theory will fail to inflate, no matter how small are the field values, because the inflaton mass is too big [3].

Either in supersymmetric or non-supersymmetric theories the slow-roll inflation is the dominant paradigm. This phase is in principle characterized by the Hubble friction, hence theories that generate enhanced friction effects are cosmologically rather motivated. It has been found that when a scalar field has derivative couplings to curvature, then it can slow-roll down even at relatively steep potentials during a (nearly) de Sitter phase. Nevertheless, not all derivative couplings to curvature are consistent, but there exist specific classes which lead to viable field theories [4]. An example of these ghost-free higher derivatives is the kinetic coupling of a scalar field to the Einstein tensor

$$\frac{1}{M_*^2} G^{mn} \partial_m \phi \partial_n \phi$$

(1)

which has given rise to a considerable amount of scientific activity [5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16, 17]. In fact there also exist even higher order consistent derivative couplings [4, 18, 19, 20, 21].

The coupling (1) has been called “Gravitationally Enhanced Friction” (GEF) mechanism [12]. The attraction of the mechanism is that it can set more general initial conditions for the inflationary phase, by relaxing the slow-roll conditions. Note that the mass scale in (1) has to be smaller than the Hubble scale during inflation because only if $M_* \ll H$ the enhanced friction effects are more influential and noticeable. This raises the question of the origin of this scale; it is rather motivating to find it among the dilaton couplings of the heterotic string [22].

Thus it would be desirable to have an embedding of the GEF mechanism in the theory of supergravity. This is in general not a trivial task; as we have mentioned, higher derivative theories may come accompanied by ghost instabilities. In earlier work [23, 24] it was understood, that in order for this coupling to be consistently realized, one has to turn to the new-minimal supergravity [25, 26, 27]. Nevertheless the coupling is still inconsistent unless the non-minimally coupled superfield has a vanishing R-charge, and is neutral under any gauge group. Thus it is not possible to endow this superfield with a conventional self-interaction in terms of superpotential or gauging. This problem can be solved by breaking supersymmetry with a Fayet-Iliopoulos term, which may induce a potential for the scalar field.

In this work we explore the modified dynamics of a non-minimally coupled superfield to curvature, and we find that inflation can be realized and described reliably in a supergravitational framework. Indeed, the scalar component of the $\Phi$ superfield is governed solely by a D-term potential and experiences a high
friction during a de Sitter phase. Moreover, due the vanishing of the superpotential, it is also expected to have negligible interactions with other fields, a fact that is supported by the Planck observational data. Therefore the superfield $\Phi$ is tailor-made for driving inflation in supergravity.

To outline, the motivation of this article is both particle theoretical and cosmological. On technical grounds we show how it is possible to introduce a potential for the non-minimally coupled field when it is coupled to supergravity (section 2), and then we show how an inflating theory driven by the supersymmetric slotheon -it has been named slotheon after [28]-, can evade some shortcomings common in conventional inflationary supergravity (section 3). In this new framework we revisit particular inflationary models, we find the new field space region where an accelerated expansion is realized (section 4) and test whether they can fit the Planck data even though they were previously excluded (section 5).

## 2 New-minimal supergravity: Derivative couplings and D-terms

The minimal theories of supergravity have a rich structure originating from the possible compensating multiplets that break the underlying superconformal theory to super-Poincaré [29, 30]. The underlying dualities among the compensating multiplets survive the gauge fixing and lead to equivalent couplings to matter [31], but break down as soon as higher derivatives are introduced. The couplings we want to study here make this duality-breakdown even more manifest, since the only known supergravity which can accommodate them in a consistent way [23, 24] is the so-called new-minimal supergravity [25, 26, 27]. An aspect of the new-minimal supergravity not encountered in the standard supergravity is the existence of a chiral symmetry. It is well known that rigid supersymmetry allows for the existence of a chiral symmetry $U(1)_R$. This $R$-symmetry becomes local and is gauged by one of the auxiliary fields of the gravitational supermultiplet.

The new-minimal supergravity [25] is the supersymmetric theory of the gravitational multiplet

$$e^a_m, \quad \psi^\alpha_m, \quad A_m, \quad B_{mn}. \quad (2)$$

The first two fields are the vierbein and its superpartner the gravitino, a spin-$\frac{3}{2}$ Rarita-Schwinger field. The last two fields are auxiliaries. The real auxiliary vector $A_m$ gauges the $U(1)_R$ chiral symmetry. The auxiliary $B_{mn}$ is a real two-form appearing only through its dual field strength $H_m$, which satisfies

$$\hat{D}^a H_a = 0 \quad (3)$$

for the supercovariant derivative $\hat{D}^a$. Indeed, the constraint (3) can be solved in terms of $B_{mn}$ as

$$H_m = -\frac{1}{3!} \epsilon_{m n r s} \left( \partial^n B^r s + \partial^r B^n s + \partial^s B^n r \right)$$

$$-\frac{1}{3!} \epsilon_{m n r s} \left( \frac{i}{8} \bar{\psi}_n \gamma^r \psi_s + \frac{i}{8} \bar{\psi}_s \gamma^n \psi_r + \frac{i}{8} \bar{\psi}_r \gamma^n \psi_s \right). \quad (4)$$

Note that all the fields of the new-minimal supergravity multiplet are gauge fields.

We will use superspace techniques to guarantee that our component form Lagrangians are supersymmetric. The interested reader may consult for example [27] where a treatment of the new-minimal superspace is given. The new minimal supergravity free Lagrangian is given by

$$\mathcal{L}_{\text{sugra}} = -M_P^2 \int d^4 \theta E \mathbf{V}_R. \quad (5)$$
Here \( V_R \) is the gauge multiplet of the R-symmetry, which (in the appropriate WZ gauge) contains the auxiliary fields in its vector component

\[
-\frac{1}{2}[\nabla_\alpha, \bar{\nabla}_{\dot{\alpha}}] V_R |_{\alpha} = A_{\alpha\dot{\alpha}} - 3H_{\alpha\dot{\alpha}}
\]

and the Ricci scalar in its highest component

\[
\frac{1}{8} \nabla^\alpha \nabla^2 \nabla_\alpha V_R |_{\alpha} = -\frac{1}{2} (R + 6H^aH_a).
\]

The superfield \( E \) is the super-determinant of new-minimal supergravity, but in general (as we also do here) one can calculate the supersymmetric Lagrangians only with the use of the F-term formula, since

\[
\int d^4\theta EX = \frac{1}{2} \int d^2\theta \mathcal{E} \left( -\frac{1}{4} \nabla^2 X \right) + c.c.
\]

and the chiral density reads

\[
\mathcal{E} = e + ie\sqrt{2}\theta\sigma^a \bar{\psi}_a - \theta^2 c\bar{\psi}_a \sigma^{ab} \tilde{\psi}_b
\]

in the chiral theta expansion. Note that \( X \) is a generic hermitian superfield with vanishing chiral weight, while its chiral projection \((-\frac{1}{4} \nabla^2 X)\) has chiral weight \( n = 1 \). The bosonic sector of Lagrangian (5) is

\[
\mathcal{L}_{\text{sugra}}^B = M_P^2 e \left( \frac{1}{2} R + 2A_aH^a - 3H_aH^a \right).
\]

For the matter sector we have a chiral multiplet, defined by

\[
\bar{\nabla}_\dot{\alpha} \Phi = 0
\]

which has bosonic components a physical complex scalar \( A \), and an auxiliary complex field \( F \), defined as

\[
\Phi |_{\alpha} = A , \quad -\frac{1}{4} \nabla^2 \Phi |_{\alpha} = F.
\]

In general, a chiral superfield in new-minimal supergravity is allowed to have an arbitrary R-charge, but we stress that our chiral superfield has a vanishing one \([23]\)

\[
n_{\Phi} = 0.
\]

The minimal kinematic Lagrangian for this multiplet is in superspace

\[
\mathcal{L}_0 = \int d^4\theta \ E \bar{\Phi} \Phi
\]

the bosonic sector of which is

\[
\mathcal{L}_0^B = A\square \bar{A} + F\bar{F} - iH^m \left( A\partial_m \bar{A} - \bar{A} \partial_m A \right).
\]

Finally, concerning our chiral superfield, it will also have a non-minimal derivative coupling with the supergravity multiplet

\[
\mathcal{L}_{M_*} = iM_*^{-2} \int d^4\theta \ E \left[ \Phi E^a \nabla_a \Phi \right] + c.c.
\]
where $E^a$ is a curvature real linear superfield ($\nabla^2 E_a = \bar{\nabla}^2 E_a = 0$) of the new-minimal supergravity. The $E^a$ superfield has bosonic components

$$E_a = H_a$$

and

$$\frac{1}{4}g^{\alpha\dot{\alpha}}[\nabla_{\alpha}, \bar{\nabla}_{\dot{\alpha}}]E_b = \frac{1}{2}(G_{ab} - g_{ab}H^cH_c - 2H_aH_b - ^*F_{ab})$$

where $G_{mn} = R_{mn} - \frac{1}{2}g_{mn}R$ is the Einstein tensor and $F_{mn} = \partial_m A_n - \partial_n A_m$ is the field strength of the supergravity auxiliary field $A_m$. It is worth mentioning that $E^a$ satisfies the superspace Bianchi identity

$$\nabla^a E_a = 0$$

which is related to the fact it contains the field strengths of all the supergravity multiplet component fields. For a discussion and derivation of the Lagrangian (16) see [23]. The bosonic sector of Lagrangian (16) is

$$L^B_M = M^{-2} \left[ G^{ab} \partial_b \bar{A} \partial_a A + 2 F \bar{F} H^a A_a - 2 F \bar{F} H^a H_a \\
+ iH^a \left( \bar{F} \partial_a F - F \partial_a \bar{F} \right) - \partial_b A \partial^b \bar{A} H_a H^a \\
+ 2 H^a \partial_a A H^b \partial_b \bar{A} - i H_c \left( \partial_b \bar{A} D^c \partial^b A - \partial_b A D^c \partial^b \bar{A} \right) \right].$$

Note that this term, although it contains higher derivatives, does not lead to ghost states or instabilities. The ghost instabilities are in fact evaded due to the vanishing chiral weight of the chiral superfield $\Phi$. On the other hand, the vanishing chiral weight forbids the self-coupling via a superpotential due to the $R$-symmetry. Thus this superfield is not allowed to have a superpotential. Moreover it is also not allowed to be gauged, since this will also give rise to ghost instabilities via inconsistent derivative couplings of the gauge fields to curvature. The only remaining option is the indirect introduction of self-interaction via a gauge kinetic function.

The gauge sector of our theory is composed of a standard $U(1)$ gauge multiplet $V$, with a $\Phi$-dependent gauge kinetic function and a Fayet-Iliopoulos term. The $U(1)$ gauge multiplet, consists of a gauge vector field $v_m$, a majorana gaugino $\lambda^a$, and a real auxiliary field $D$. In particular, the definition of the bosonic components of the vector is

$$-\frac{1}{2}[\nabla_{\alpha}, \bar{\nabla}_{\dot{\alpha}}]V = v_{a\dot{a}}, \frac{1}{8} \nabla^a \nabla^2 \nabla_a V = D.$$  

Note that due to the structure of new-minimal supergravity, a FI term is in general allowed (even the superspace Lagrangian of pure new-minimal supergravity is a FI term). Thus we have in superspace

$$L_g = \frac{1}{4} \int d^2 \theta \mathcal{E} f(\Phi) W^2(V) + c.c. + 2\xi \int d^4 \theta E V$$

with

$$W_a(V) = -\frac{1}{4} \nabla^2 \nabla_a V$$

where $f(\Phi)$ is a holomorphic function of the chiral superfield $\Phi$ and $\xi$ is the Fayet-Iliopoulos parameter of mass dimension two. The bosonic sector of (22) reads

$$e^{-1} L_g^B = -\frac{1}{4} \text{Re} f(A) F^{mn} F_{mn} + \frac{1}{4} \text{Im} f(A) F^{mn} * F_{mn} + \frac{1}{2} \text{Re} f(A) D^2 + \xi D - 2\xi v_a H^a$$

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where

\[ F_{mn} = \partial_m v_n - \partial_n v_m. \]  

(25)

We stress that the scalar \( A \) is not charged under this \( U(1) \), thus there is no restriction in the form of \( f(A) \) apart from holomorphicity.

The total Lagrangian we are interested in is

\[ \mathcal{L}_{\text{total}} = M_P^2 \mathcal{L}_{\text{sugra}} + \mathcal{L}_g + \mathcal{L}_0 + \mathcal{L}_M, \]

(26)

which reads

\[
e^{-1} \mathcal{L}_{\text{total}} = M_P^2 \left[ \frac{1}{2} R + 2 V^a H_a - 3 H_a H_a \right] + A \Box \bar{A} + F \bar{F} \\
+ M^{-2} \left[ G^{ab} \partial_b \bar{A} \partial_a A - 2 F \bar{F} H^a H_a - \partial_b A \partial^b \bar{A} H_a H_a + 2 H^a \partial_a A H^b \partial_b \bar{A} \right] \\
- \frac{1}{4} \text{Re} f(A) F^{mn} F_{mn} + \frac{1}{4} \text{Im} f(A) F^{mn} \ast F_{mn} \\
+ \frac{1}{2} \text{Re} f(A) D^2 + \xi D \]

(27)

where we have redefined the auxiliary field \( A^a \) to \( V^a \)

\[
V^a = A^a \left( 1 + \frac{1}{M_P^2} M_s^{-2} F \bar{F} \right) - \frac{1}{M_P^2} \xi v^a \\
+ \frac{1}{2 M_P^2} \left( i \bar{A} \partial^a A - i A \partial^a \bar{A} - i M_s^{-2} F \partial^a \bar{F} + i M_s^{-2} \bar{F} \partial^a F \right) \\
+ \frac{1}{2 M_P^2} \left( i M_s^{-2} \partial_b A D^a \partial^b \bar{A} - i M_s^{-2} \partial_b \bar{A} D^a \partial^b A \right). \]

(28)

Lagrangian (27) contains four auxiliary fields. First, by solving the equations of motion for the supergravity auxiliary fields we find that the vector \( H_m \) vanishes and \( V_m \) reduces to a pure gauge. In fact since \( H_m \) is the dual field-strength of \( B_{mn} \) here both auxiliary fields of the new-minimal supergravity are pure gauge on-shell. For the auxiliary field \( F \) it is easy to see that it will also vanish on-shell. Finally, by solving the equations of motion for the auxiliary field \( D \) of the gauge multiplet we find

\[ D = - \frac{\xi}{\text{Re} f(A)}. \]

(29)

After plugging back our results, we have the following on-shell form for (27)

\[
e^{-1} \mathcal{L}_{\text{total}} = \frac{M_P^2}{2} R + A \Box \bar{A} + M_s^{-2} G^{ab} \partial_a \bar{A} \partial_b A - \frac{1}{2} \frac{\xi^2}{\text{Re} f(A)} \\
- \frac{1}{4} \text{Re} f(A) F^{mn} F_{mn} + \frac{1}{4} \text{Im} f(A) F^{mn} \ast F_{mn}. \]

(30)

Note that this Lagrangian (30) does not contain ghost states or instabilities.

The Fayet-Iliopoulos term inside (24) breaks supersymmetry, and combining it with the gauge kinetic function has the effect of introducing a scalar potential of the following form

\[ V = \frac{1}{2} \frac{\xi^2}{\text{Re} f(A)}. \]

(31)
This is the D-term potential. It is expected that only in the case of broken supersymmetry one can have a potential for the slotheon $\lambda$ field. This stems from the fact that the chiral $U(1)_R$ symmetry of new-minimal supergravity forbids a potential for this field, due to its vanishing chiral weight. The advantage of a Fayet-Iliopoulos term is that it breaks supersymmetry spontaneously.

3 Application to inflation

3.1 A pure D-term inflation

In the standard supergravity the scalar potential of chiral superfields transforming in some representation of a gauge group has the following form:

$$V = e^{K/M_P^2} \left[ F_i (K^{-1})^j_i F^j - 3 |W|^2 M_P^2 \right] + \frac{g^2}{2} \frac{1}{Re f_{ab}} D^a D^b$$

(32)

where $F^i = W_i + K^i W/M_P^2$ and $D^\alpha = K^i (T^a)^j_i z_j + \xi^a$. The upper (lower) index $i$ denotes derivatives with respect to the $\phi_i$ ($\phi^*_i$) field. The slow-roll conditions imply

$$\epsilon \ll 1 \Rightarrow \frac{K_{\phi}}{M_P} + \ldots \ll 1$$

(33)

$$\eta \ll 1 \Rightarrow 3K_{\phi}\phi H^2 + \ldots \ll H^2.$$  

(34)

Here the subscript $\phi$ denotes a derivative with respect to the inflaton. The inflaton vacuum energy dominates the energy density of the universe and the relation $H^2 = V/(3M_P^2)$ has been used in the second condition $\eta \ll 1$. In the low energy minimum the Kähler metric should be normalized to one and it is not expected to be suppressed during inflation. Therefore F-type inflation in supergravity theories is hard to be realized unless the Kähler and the superpotential have a special form or accidental cancellations take place \[3, 1, 35\].

A resolution to this $\eta$-problem in generic supergravity theories can be given by a symmetry that suppresses the F-term part of the scalar potential. Thereby the potential is naturally dominated by a Fayet-Iliopoulos D-term which exists for $U(1)$ gauge groups. The R-symmetry of the theory forbids the superpotential interactions for the $A$ field non-minimally coupled to the $G^mn$ tensor. The spontaneous breaking of supersymmetry during inflation may introduce interactions however these will be generated radiatively and should not affect the tree level D-term inflationary potential. The D-term potential domination together with the enhanced friction features strongly motivates the study of this higher derivative theory to inflationary applications.

3.2 Expanding the allowed initial conditions for inflation

The complex scalar field $A$ is governed by the scalar potential generated by the Fayet-Iliopoulos supersymmetry breaking and has the form \[35\]

$$V(\phi, \beta) = \frac{1}{2} \frac{\xi^2}{Re f(\phi, \beta)}$$

(35)

\[1\] This formula is common for the old-minimal supergravity \[32\]. Nevertheless, a general supergravity-matter system in the new-minimal framework can be recast in this form after appropriate redefinitions \[27\].

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where \( \phi = \text{Re}\{A\} \) and \( \beta = \text{Im}\{A\} \). In our context the gauge kinetic function \( \text{Re} f(A) \) is arbitrary and in principle contains non-renormalizable terms. In most of the models, again, one finds that the \( |A| \) is of order \( M_P \) or larger, a fact that makes the non-renormalizable terms difficult to control similarly to the higher order terms in the \( \mathcal{K} \) and \( \mathcal{W} \) potential. Here, we will approximate \( f(A) \) by a polynomial expansion or ascribe to it forms suggested by microscopic theories as the string theory. The fact that the complex \( A \)-field has no superpotential also implies that potential \((35)\) does not receive radiative corrections from superpotential interactions.

In a FLRW background, neglecting spatial gradients, the Friedmann equation and the equation of motion for the \( \phi \) (or the \( \beta \) field) are

\[
H^2 = \frac{1}{3M_P^2} \left[ \frac{\dot{\phi}^2}{2} \left(1 + 9M_*^{-2}H^2\right) + V(\phi) \right], \quad \partial_t \left[ a^3\dot{\phi} \left(1 + 3M_*^{-2}H^2\right) \right] = -a^3V_\phi. \tag{36}
\]

Let us first demonstrate the advantages of the kinetic coupling to the inflationary applications. We assume here that the \( \phi \) is the single inflating field and we consider a full polynomial potential without symmetry suppressed terms, actually non-renormalizable terms are naturally present in supergravity theories. Thus,

\[
V(\phi) = V_0 + m_1^3\phi + \frac{1}{2}m_2^2\phi^2 + \frac{1}{3}m_3^3\phi^3 + \frac{1}{4}\lambda_4\phi^4 + \sum_{n=5}^\infty \lambda_nM_P^{4-n}\phi^n. \tag{37}
\]

In the large field models of inflation the inflaton field has a value of order the Planck mass, \( M_P \). The above general potential cannot serve as large field inflationary model for the non-renormalizable terms, which may be unreasonable to be neglected, spoil the flatness of the potential.

The non-minimal coupling of the kinetic term of the scalar field with the Einstein tensor \( G_{mn} \)

\[
\mathcal{L} = -\frac{1}{2}\sqrt{-g} \left(g^{mn} - M_*^{-2}G^{mn}\right) \partial_m \phi \partial_n \phi \tag{38}
\]

during a de Sitter phase takes the simple form \( M_*^{-2}G^{mn} = -3M_*^{-2}H^2 g^{mn} \). For \( HM_*^{-1} \gg 1 \) the kinetic coupling implies that the canonically normalized scalar field is \( \tilde{\phi} = 3HM_*^{-1}\phi \). This rescaling recasts the polynomial potential \((37)\) in terms of the canonically normalized inflaton \( \tilde{\phi} \) to the form

\[
V(\tilde{\phi}) = V_0 + m_1^3\frac{\tilde{\phi}}{3HM_*^{-1}} + \frac{1}{2}m_2^2\left(\frac{\tilde{\phi}}{3HM_*^{-1}}\right)^2 + \frac{1}{3}m_3^3\left(\frac{\tilde{\phi}}{3HM_*^{-1}}\right)^3 + \frac{1}{4}\lambda_4\left(\frac{\tilde{\phi}}{3HM_*^{-1}}\right)^4 + \\
+ \sum_{n=5}^\infty \lambda_nM_P^{4-n}\left(\frac{\tilde{\phi}}{3HM_*^{-1}}\right)^n. \tag{39}
\]

The non-renormalizable terms \( \sum_{n=4}^\infty \lambda_nM_P^n \left(\tilde{\phi} \times (3HM_*^{-1}M_P)^{-1}\right)^n \) are suppressed by the “enhanced” mass scale \( 3HM_*^{-1}M_P \). The slow roll parameters require \( \tilde{\phi} > M_P \) and, hence, these higher order terms can be neglected and sufficient inflation can take place given that

\[
M_P < \tilde{\phi} \ll M_P(HM_*^{-1}). \tag{40}
\]

In terms of the field \( \phi \), which has non-canonical kinetic term, the above field-space region translates into

\[
\frac{M_P}{HM_*^{-1}} < \phi \ll M_P. \tag{41}
\]
This finding is of central importance since we work in a supersymmetric context. Even though we suggest a D-term inflation without a superpotential the generation of the inflationary potential is in principle not protected by any symmetry and in the most general case we cannot forbid the higher order terms.

In this work we focus on inflationary potentials of the form \( \text{[37]} \). We consider our theory as an effective one valid below some ultra-violet cut-off that we generally identify with the \( M_P \). The field-space region \( \text{[11]} \) allows inflation to be realized in general form of potentials and reliable conclusions in this context can be derived. It can be said that the kinetic coupling theory is tailor-made for realizing an inflationary phase.

From a different perspective, if there is an internal symmetry that forbids the non-renormalizable terms and thereby suppresses the coefficients \( \lambda_n \) for \( n \geq 5 \) then inflation can be implemented in a much larger field-space region, than in the conventional (GR limit) large field inflationary models, that reads:

\[
\phi > \frac{M_P}{H M_*}. \tag{42}
\]

### 4 Introducing D-term inflationary potentials

We will attempt to capture some of the characteristics of the kinetic coupling in inflationary applications by considering some representative examples of inflationary potentials. We will concentrate on single field inflation models assuming that one of the two fields is stabilized in the vacuum.

According to the equations \( \text{[36]} \) and for \( H M_*^{-1} \gg 1 \) the slow-roll parameters of General Relativity (GR) \( \epsilon \equiv M_P (V'/V)^2/2 \) and \( \eta \equiv M_P^2 V''/V \) are recast into

\[
\tilde{\epsilon} \approx \frac{\epsilon}{3 H^2 M_*^{-2}} \quad \tilde{\eta} \approx \frac{\eta}{3 H^2 M_*^{-2}}. \tag{43}
\]

The requirement \( \tilde{\epsilon}, |\tilde{\eta}| < 1 \) yields that the field space region where slow-roll inflation is realized is rather increased. We will illustrate this below by considering different forms for the gauge kinetic function and thereby various types of potentials.

**Hill-top models.** Let us first assume that

\[
f(A) = \frac{\xi^2}{2 V_0} \sum_n \lambda_n \left( \frac{A}{M_P} \right)^n \tag{44}
\]

where \( \lambda_n \) are real coefficients and we constrain the field to subplanckian values \( |A| < M_P \). Then \( f(A) = (\xi^2/2 V_0) (1 + \lambda_1 A/M_P + \lambda_2 A^2/M_P^2 + ... ) \) and the Re \( f(A) \) is maximized for \( |A| \to \infty \) for \( \lambda_n > 0 \). The scalar potential reads \( V(\phi, \beta) = V_0 \left( 1 - \lambda_1 \phi/M_P - (\lambda_2 - \lambda_2^1) \phi^2/M_P^2 + \lambda_2 \beta^2/M_P^2 + ... \right) \) where the ellipsis correspond to negligible terms. The above potential includes two scalars that have a non-minimal derivative coupling. It may correspond to a multi field inflation. Here we consider that \( \langle \beta \rangle = 0 \) which is a \( \beta \)-direction minimum. The inflating scalar is the \( \phi \) with potential

\[
V = V_0 \left( 1 - \lambda_1 \frac{\phi}{M_P} - \lambda' \frac{\phi^2}{M_P^2} + ... \right) \tag{45}
\]

where \( \lambda' \equiv \lambda_2 - \lambda_2^1 \), which is valid for \( \phi \ll M_P \) with the \( \phi \) field rolling down to larger values. If the linear to \( \phi \) term is forbidden by symmetries then we obtain the quadratic “hill-top” model for inflation \( V = V_0 \left( 1 - \lambda' \phi^2/M_P^2 + ... \right) \) otherwise the linear term dominates for small values for the field and the potential corresponds to a linear “hill-top” model. The form of the \( \text{[45]} \) potential is an approximation
about the origin, and one can see that it is bounded from below when the full terms are taken into account.

The slow-roll conditions $\tilde{\epsilon}, |\tilde{\eta}| \ll 1$ yield the requirements for inflation

\[
\left\{ \begin{array}{l}

\tilde{\epsilon} \approx \frac{M_P^2 \lambda^2}{(2V_0 M_*^{-2})} < 1 \\

\tilde{\eta} = 0
\end{array} \right. \quad \text{and} \quad \phi \ll M_P
\]  

(46)

for the linear potential, and

\[
\left\{ \begin{array}{l}

\tilde{\epsilon} \approx 2\lambda^2 \phi^2/(V_0 M_*^{-2}) < 1 \\

\tilde{\eta} \approx |2\lambda^2 M_P^2/(2V_0 M_*^{-2})| < 1
\end{array} \right. \quad \Rightarrow \quad \phi < \frac{V_0^{1/2} M_*^{-1}}{\sqrt{2\lambda^2}} \ll M_P
\]  

(47)

for the quadratic hill-top model, i.e. when the odd terms are absent, $\lambda_1 = 0$. Obviously when $\lambda_1 = \mathcal{O}(1)$ in the linear case and $\lambda = \mathcal{O}(1)$ in the quadratic one, inflation cannot be realized without the kinetic coupling. Indeed the slow-roll conditions $\tilde{\epsilon}, |\tilde{\eta}| < 1$ are violated and one usually asks for a mass scale larger than the $M_P$ to suppress the powers of the $\phi$ field in the expansion $[45]$. The kinetic coupling addresses this problem and allows sufficient inflation to take place.

The hill-top models can be compatible with the CMB data $[33]$ only as a large field model and for $M_P/\lambda' \gtrsim 9M_P$. Obviously, when $\lambda' = \mathcal{O}(1)$ the model is excluded by Planck data. Here, thanks to the kinetic coupling it is possible the field $\phi$ to drive inflation for sub-Planckian values and for a coefficient $\lambda' = \mathcal{O}(1)$. Indeed, the canonically normalized scalar field during inflation, $\tilde{\phi}$, has a potential $\mathcal{V} \approx V_0 (1 - \lambda' \phi^2 / M_P^2 + ...)$ and the effective coupling is the $\lambda'$ where $\lambda' = \lambda'/(3HM_*^{-1})^2 \ll \lambda'$.

**Power law potentials (chaotic) models.** If the gauge kinetic function is of the form

\[
\mathcal{f}(A) = \frac{\xi^2}{2V_0} \left[ \sum_n \lambda_n \left( \frac{A}{M_P} \right)^n \right]^{-1}
\]  

(48)

then for $|A| < M_P$ $\mathcal{f}(A) = (\xi^2 / 2V_0)(1 + \lambda_1 A / M_P + \lambda_2 A^2 / M_P^2 + ...)^{-1}$. The Re $\mathcal{f}(A)$ for $\lambda_n > 0$ is maximized when $|A| \to 0$. The scalar potential is $\mathcal{V}(\phi, \beta) \approx V_0 \left( 1 + \lambda_1 \phi / M_P + [\lambda_2 \phi^2 - \beta^2(\lambda_2 - \lambda_2^2)] / M_P^2 + ... \right)$ and the chaotic type large field models potential is obtained. Assuming that $\lambda_1, \lambda_2 > 0$ then the potential for sub-Planckian values for the fields reads approximately $^2$

\[
\mathcal{V}(\phi, \beta) \approx V_0 \left( 1 + \frac{\lambda_1}{M_P} \phi + \frac{\lambda_2^2}{M_P^2} \beta^2 + ... \right).
\]  

(49)

This is a large field model potential in the sense that the fields roll towards the origin during inflation. However, also here, this potential approximates the dynamics of the fields for $|A| < M_P$ where our theory is under control. When $\beta < \phi$ the linear to $\phi$ potential term dominates which can inflate the universe. On the other hand, if $\phi < \lambda_1 \beta^2 / M_P$ then a quadratic to $\beta$-field potential may drive inflation.

The inflating field slow-rolls when the two conditions $\tilde{\epsilon}, |\tilde{\eta}| \ll 1$ are satisfied. After naming $m_1^3 \equiv \lambda_1 V_0 / M_P$ and $m_2^3 \equiv 2\lambda_2^2 V_0 / M_P$ we take the slow-roll conditions for the limiting cases of a linear and a quadratic potential. These read respectively

\[
\left\{ \begin{array}{l}

\tilde{\epsilon} = M_P^2 / (2m_1^3 \phi^3 M_*^{-2}) < 1 \\

\tilde{\eta} = 0
\end{array} \right. \quad \Rightarrow \quad \left( \frac{M_P^4}{2m_1^3 M_*^{-2}} \right)^{1/3} < \phi \ll M_P
\]  

(50)

$^2$Note that the same form of the potential can be obtained from the gauge kinetic function $[44]$ for $|A| < M_P$, $\lambda_1 < 0$ and $\lambda_2 - \lambda_2^2 < 0$. 

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for the linear potential, and
\[
\begin{align*}
\{ \dot{\epsilon} &= 4M_P^4/(m_2^2\beta^4M_*^{-2}) < 1 \\
\eta &= \dot{\epsilon} < 1 \}
\end{align*}
\Rightarrow \frac{\sqrt{2}M_P}{(m_2^2M_*^{-1})^{1/2}} < \beta \ll M_P
\tag{51}
\]
for a quadratic potential where the $\phi$ field is about the origin $\phi \approx 0$. We note here that the Planck data analysis [33] favours models with weak or vanishing tensor-to-scalar ratio and negative curvature $V'' < 0$. In particular the quadratic potential is disfavoured. We will show in the next section that due to the kinetic coupling these well motivated models can still fit the Planck data.

**Exponential potentials.** If we now assume that the gauge kinetic function is of the form
\[
f(A) = \frac{\xi^2}{2V_0} e^{\lambda A/M_P}
\tag{52}
\]
then we directly get an exponential type potential
\[
V = V_0 \frac{1}{\cos(\lambda/\sqrt{MP})} e^{-\lambda\phi/M_P}.
\tag{53}
\]
The form of the function $1/\cos x$ suggests that the $\beta$-dependent part of the potential will be stabilized to values $(1/\cos(\lambda\beta/M_P)) = 1$ and the $\phi$ will be the inflating field. An important reason for picking the gauge kinetic function [52] is that it is reminiscent of the dilaton coupling of string theory. Slow-roll inflation takes place for
\[
\begin{align*}
\{ \dot{\epsilon} &= (\lambda^2/2)/(3H^2M_*^{-2}) < 1 \\
\eta &= 2\dot{\epsilon} < 1 \}
\end{align*}
\Rightarrow \phi < \frac{M_P}{\lambda} \ln \left( \frac{2V_0M_*^{-2}}{M_P^2\lambda^2} \right).
\tag{54}
\]
For this type of potentials an inflationary phase can be realized for a wider range of parameters as well. Namely, in GR limit, inflation driven by an exponential potential is impossible for $\lambda = O(1)$. It has to be $\lambda < 1$ which implies, after absorbing $\lambda$ to the mass scale in the denominator, that the field $\phi$ has to be suppressed by a super-Planckian value. Here thanks to the kinetic coupling inflation is possible even for $\lambda \gtrsim 1$.

**Inverse power law potentials.** Finally let us consider that
\[
f(A) = \frac{\xi^2}{2V_0} \frac{A^n}{\lambda M_P^n}.
\tag{55}
\]
Using the relation Re$\{A^n\} = \phi^n \cos(n\theta)/(\cos \theta)^n$ we see that the field $\theta$, the phase of the complex field $A = \rho e^{i\theta}$, is stabilized at $\theta = \kappa \pi$ and so Re$\{A^n\} = \phi^n$. The potential reads
\[
V = V_0 \frac{1}{\lambda} \frac{M_P^n}{\phi^n},
\tag{56}
\]
and slow-roll inflation takes place for
\[
\begin{align*}
\{ \dot{\epsilon} &= (M_P^2n^2)/(2\phi^2 \times 3H^2M_*^{-2}) < 1 \\
\eta &= 2(1+n^{-1})\dot{\epsilon} < 1 \}
\end{align*}
\Rightarrow \phi^{n-2} < 2M_P^{-4} \frac{V_0M_*^{-2}}{\lambda n^2}.
\tag{57}
\]
Again here, for the inverse power law potential, inflation can be achieved in a larger part of the $\phi$-field space. In particular sub-Planckian values for the field can drive inflation whereas in the GR limit this is ruled out since there the slow-roll conditions require $\phi \gtrsim M_P$. 

11
5 Enhanced friction supergravity inflationary models and the CMB data

The Planck satellite measurement of the temperature anisotropies in Cosmic Microwave Background (CMB) is the ruler for model selection. In order to make contact between the theory and observation the spectra of scalar and tensor perturbations have to be estimated [1, 34]. The density perturbations $\delta \rho$ of the inflaton are encoded in the variable $\zeta = \delta \rho / (\rho + p)$ which is conserved on large scales in the absence of entropy perturbations and can be directly related to the cosmic microwave background temperature fluctuations. The power spectrum of the $\zeta$ variable in first order in the slow roll parameter $\tilde{\epsilon}$ reads [12]

$$P_\zeta \equiv \frac{k^3}{2\pi^2} |\zeta_k|^2 \approx \frac{H^2}{8\pi^2 \epsilon \zeta M_P^2}$$

(58)

where the sound speed squared having a dependence $c_s^2 \propto H^2 \tilde{\epsilon}$, is subluminal and modifies the spectral tilt dependence on the slow roll parameters. The result is

$$n_s - 1 \equiv \frac{d \ln P_\zeta}{d \ln k} \bigg|_{\epsilon, k = a H} \approx -8 \tilde{\epsilon} + 2 \eta$$

(59)

counter to the well known GR limit formula $n_s - 1 = 2 \eta - 6 \epsilon$. The ratio of the tensor to scalar amplitudes, $r \equiv P_{\eta}(k_s) / P_\zeta(k_s)$, has the conventional GR dependence on the slow-roll parameter $\tilde{\epsilon}$ at lowest order: $r = 16 \tilde{\epsilon}$. However the new relation (59) allows for larger values for the $\eta$ slow-roll parameter. Namely it is

$$r = 2(1 - n_s) + 4 \eta$$

(60)

instead of $r = (8 / 3)(1 - n_s) + (16 / 3)\eta$. Hence, given that $1 - n_s \approx 0.04$ and $r < 0.11$ [33], positive values for the $\eta$ can be accommodated, which correspond to potentials with $V'' \geq 0$. The number of e-folds, $N \equiv \int H dt$, for $H^2 M_s^{-2} \gg 1$ is given by the expression

$$N(\phi) = \frac{1}{M_P^2} \int_{\phi_f}^{\phi_i} \left(1 + 3 H^2 M_s^{-2} \right) \frac{V}{V'} d\phi \approx \frac{1}{M_P^2} \int_{\phi_f}^{\phi_i} M_s^{-2} \frac{V^2}{V'} d\phi$$

(61)

The Planck collaboration estimated the spectral index $n_s$ from the observational data (Planck and WMAP) to be [33] $n_s = 0.9603 \pm 0.0073$ and the upper bound on the tensor to scalar ratio at $r < 0.11$. This constraint on $r$ corresponds to an upper bound on the energy scale of inflation $H(\phi_s) / M_P \leq 3.7 \times 10^{-5}$ which implies that $\epsilon(\phi_s) \equiv \tilde{\epsilon}_s < 0.008$. The $\phi_s$ denotes the field value during inflation that the pivot scale $k_s = 0.002 \text{Mpc}^{-1}$ [33] exited the Hubble radius (not to be confused with the star index at the mass scale $M_s$ of the non-minimal derivative coupling).

Let us now examine the supergravity D-term inflationary models of the previous section in the light of the Planck data. The key ingredient is that the inflaton field is characterized by the non-minimal derivative coupling to the Einstein tensor.

The hill-top model (45) when the linear to $\phi$ term dominates the potential yields a spectral index $1 - n_s = 8 \tilde{\epsilon} = 4 \lambda_2^2 M_p^2 M_s^2 / V_0$ which is related to the number of e-folds by the expression $1 - n_s \approx 4 \lambda_2 (\phi_f - \phi_s) / (N(\phi_s) M_P)$ where $\phi_f$ is the field value at the end of inflation. For $N(\phi_s) \sim 50$ and $\Delta \phi \ll M_p$ this model predicts a spectral index $n_s < 0.04$ and it is in tension with the Planck data. On the other hand the quadratic hill-top model, which is derived when the gauge kinetic function $f(A)$ is characterized by a reflection symmetry, yields a spectral index $1 - n_s \approx -2 \eta = N(\phi_s)^{-1} \ln(\phi_f / \phi_s)$ and negligible tensor to scalar ratio. The higher derivative quadratic hill-top model is in agreement
with the Planck data without the need for super-Planckian values neither for the inflaton field nor for the suppression scale $M_P/\lambda$. We also note that the quadratic hill-top model is directly related to the natural inflation models; these have been shown, in the same context but without supergravity, to be UV protected $[9, 12]$.

The power-law models $[19]$ in the linear limit yield a spectral index $1 - n_s = 8\tilde{\epsilon} = 4N(\phi_*)^{-1}/3$ and a ratio $r$ which are in better agreement with the Planck data than the linear power-law model in the GR limit. The quadratic power-law model, obtained for $\phi \approx 0$, and the $\beta$ field displaced from the origin yields $1 - n_s = 6\tilde{\epsilon} = 3N(\phi_*)^{-1}/2$ which is well within the allowed parameter region of the Planck Data. This result is of great importance because the quadratic potential appears generic since it can be interpreted as the lowest order term in a Taylor expansion about the origin. Despite its simplicity and rationalness the quadratic potential in the GR limit does not provide a good fit to the Planck data for $N(\phi_*) \lesssim 60$ e-folds. The kinetic coupling saves this model predicting $\tilde{\epsilon}_* = 0.0067$, $N(\phi_*) \approx 38$ and $r = 0.1$ with sub-Planckian variation for the value of the inflaton field.

The exponential potential and inverse power-law potentials are excluded by the CMB data even when we consider a non-minimal derivative coupling between the Einstein tensor and the scalar field. In the case of an exponential potential $[53]$ it is $\tilde{\eta} = 2\tilde{\epsilon}$ and so $1 - n_s = 6\tilde{\epsilon} - 2\tilde{\eta} = 4\tilde{\epsilon}$ which gives $\tilde{\epsilon}_* = 0.01$. Although this value approaches the observational bound $[33]$ the model still lies outside the preferred region indicated by the Planck collaboration. In the case of the inverse power-law models the spectral index is given by the modified expression $1 - n_s = 4\tilde{\epsilon}(1 - n^{-1})$ and is also ruled out by the CMB analysis. However, the kinetic coupling operates like an enhanced friction and inflation takes place more generically than in the GR limit. Since inflation is primarily introduced to address (or better ameliorate) the homogeneity, isotropy and flatness problem we can say that these two last models are still motivated candidates for inflation in the context of supergravity kinetic coupling. Afterwards, in order to seed the large-scale structure formation in the universe, a mechanism like the curvaton or the modulated reheating may take place in the post-inflationary universe.

The non-minimal derivative coupling of the scalar field $A = \phi + i\beta$ with the Einstein tensor in supergravity renders the scalar $A$ a compelling inflaton candidate due to both the enhanced friction effect and the symmetry suppression of the F-terms. Despite these advantages the absence of superpotential interactions may be problematic for a sufficient reheating of the universe. Although it is not the scope of this work to fully describe the reheating issues there are some remarks about the entropy production from the $A$-decay in order. Firstly, the inflating scalar $A$ has a coupling with the gauge field strength $[30]$. Secondly, the supersymmetry breaking, due to the non-zero vacuum energy in the early universe, may allow interactions beyond the tree level between the scalar and other fields. Finally, we note that there are also mechanisms like the gravitational particle production $[36, 37]$ invented for such questionable situations.

### 6 Conclusions

In the present paper we have examined the implementation of an inflationary phase by the scalar component of a chiral superfield in a supergravity context. The particular characteristic of this scalar is that it is non-minimally coupled to the Einstein tensor $G^{mn}$, and the slow-roll conditions can be satisfied more generically. The potential is introduced via a Fayet-Iliopoulos D-term since the R-symmetry of the theory excludes the introduction of a superpotential, and its gauging is also forbidden due to stability issues. Even though this is a higher derivative theory it does not give rise to ghost instabilities.

The model we propose is a pure D-term inflation with a gravitational enhanced slow-roll for the inflaton. These two features have important implications for the inflationary dynamics. Firstly, the accelerated expansion can be realized for sub-Planckian values for the inflaton field as well as for sub-
Planckian parameter scales. Sub-Planckian values for the inflaton field are welcome because the non-renormalizable terms are suppressed. This fact together with the absence of the F-terms render this model free from the notorious $\eta$-problem of supergravity. Secondly, the field space region where the slow-roll conditions are satisfied is increased. Hence an inflationary period is realized for more generic initial conditions. Thirdly, the relation between the spectral index of the scalar perturbations and the slow-roll parameters is modified due to the corrected sound speed of the scalar perturbations. Also a sufficient number of e-folds can be easier achieved. These imply that some inflationary models may provide a better fit to the Planck data or even render some excluded inflationary models observationally viable. For example, here, the discrepancy between the predictions of the quadratic potential, $V \propto \phi^2$ and the Planck data can be reconciled.

Concluding, some comments are in order. Throughout this work we have considered potentials characterized by $M_P$ suppression scale. Moreover, we have not taken into account multi-field implications to the dynamics of the system which may give rise to effects that deserve a further study. Finally, we would like to mention that the origin of this new scale $M_*$ has remained unspecified.

We believe that this work has offered some different insight into how inflation might work in a supergravity framework.

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