Gauge inflation by kinetic coupled gravity

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Abstract

Recently, a new class of inflationary models, so called gauge-flation or non-Abelian gauge field inflation has been introduced where the slow-roll inflation is driven by a non-Abelian gauge field $A$ with the field strength $F$. This class of models are based on a gauge field theory having $F^2$ and $F^4$ terms with a non-Abelian gauge group minimally coupled to gravity. Here, we present a new class of such inflationary models based on a gauge field theory having only $F^2$ term with non-Abelian gauge fields non-minimally coupled to gravity. The non-minimal coupling is set up by introducing the Einstein tensor besides the metric tensor within the $F^2$ term, which is called kinetic coupled gravity.

1 Introduction

The idea of inflation is a well known scenario to overcome the problems of standard cosmology [1, 2]. It tells us that the early universe has experienced an inflationary expansion phase in a very short period of time. This scenario is also very successful in reproducing the current cosmological data through the ΛCDM model [3]. One of the key ingredients in inflation is the existence of a scalar field which is subjected to the slow-roll approximation, where the kinetic energy of the scalar field remains sufficiently small compared to its potential energy. The universe experiences an inflationary expansion in the slow-roll regime, and once the kinetic energy becomes comparable to its potential energy the inflation is ended. The scalar field gets more and more kinetic energy and reaches the minimum of potential, afterwards the scalar field starts fast frictional oscillations around this minimum and transfer its kinetic energy into the matter or radiation to establish the reheating phase. Inflationary models typically benefit of single or multi-scalar field theories. However, it is very appealing to model the inflationary scenario within the particle physics context. This is because the inflation is supposed to happen at very early universe where use of high energy physics is inevitable.

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Other than scalar fields, the vector gauge fields are among the most suitable candidates to set up inflation in the context of high energy physics [5]. Nevertheless, a successful vector inflation is not possible in a *gauge invariant* manner [6].

To overcome this problem, a new class of vector inflation models is proposed in the framework of gauge field theories, so called *Gauge-flation* [7]. To remedy the incompatibility between vector gauge fields in high energy physics and requirement of isotropy in cosmology, three gauge fields are introduced with rotational symmetry resulting from $SU(2)$ non-Abelian gauge transformations. The global part of this gauge symmetry is then identified with the rotational symmetry in 3d space and so the isotropy of space in the presence of this vector gauge field is preserved.

Recently, in order to justify the current acceleration of universe, a considerable amount of activity has been focused on the kinetic coupled gravity theories [8, 9, 10, 11]. In these theories, a non-canonical kinetic term or a non-minimal coupling to the metric is responsible for the physics of acceleration. In a very recent work, the authors have introduced a new scenario of scalar field inflation in the context of kinetic coupled gravity [12]. In the present work, we aim to generalize the scenario of scalar field inflation to gauge field inflation in the context of kinetic coupled gravity. In this regard, we combine the scenario of gauge-flation [7] with the idea of kinetic coupled gravity. In explicit words, we replace $F_2$ and $F_4$ terms minimally coupled to gravity by $F^2$ term non-minimally coupled to gravity. The non-minimal coupling is established by considering the Einstein tensor besides the metric tensor within the $F^2$ term, which is called kinetic coupled gravity.

## 2 Gauge Inflation by non-Abelian $F^2$ and $F^4$ terms minimally coupled to gravity

In this section we briefly introduce the idea of gauge inflation extensively discussed in [7]. Let us consider a 4-dimensional $SU(2)$ gauge field $A^a_\mu$, where $a, b, ...$ and $\mu, \nu, ...$ are used respectively for the indices of the gauge algebra and the space-time coordinates. Gauge and Lorentz invariant Lagrangians $\mathcal{L}(F^a_\mu, g_{\mu\nu})$ are constructed out of the metric tensor $g_{\mu\nu}$ and the field strength tensor

$$F^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu - g\epsilon^{a}_{bc}A^b_\mu A^c_\nu,$$

where $\epsilon_{abc}$ is the totally antisymmetric tensor. We take flat ($k = 0$) FRW background metric

$$ds^2 = -dt^2 + a(t)^2\delta_{ij}dx^idx^j,$$

where indices $i, j, ...$ label the spatial coordinates. The effective inflaton field is introduced in the following way. If we consider an Abelian gauge field and turn on a vector gauge field in the background, then the rotational symmetry of the flat FRW geometry will be broken down, unless only we turn on the temporal component which should be only time dependent in order to preserve homogeneity. Hence, this gauge field becomes a pure gauge with vanishing field strength which is not physically viable. However, one may consider non-Abelian gauge fields with gauge group indices. It is obvious that two gauge fields which
are related by a local gauge transformation are physically equivalent. In the coordinate system of standard spatially flat FRW metric subject to the homogeneity we may turn on a typical vector gauge field with spatial components \( A_i = A_i(t) \). The time component \( A_0 \) may always be set to zero in the temporal gauge. This choice fixes the gauge transformations up to time independent gauge transformations. On the other hand, since we have already ignored the space dependence of the gauge fields, it turns out that the gauge freedom is fixed up to space-time independent global gauge transformations. Therefore, up to global gauge transformations, the background gauge field is given by

\[
A_0 = 0, \quad A_i = A_i(t). \tag{2.3}
\]

The global gauge transformations can remedy the problem of rotational symmetry breaking for non-Abelian gauge fields. Two non-Abelian gauge fields related by a gauge transformation are physically equivalent as follows

\[
(A_i)_G = U^{-1}A_iU. \tag{2.4}
\]

Moreover, the global spatial rotations are defined by

\[
(A_i)_R = R_{ij}A_j. \tag{2.5}
\]

Now, if one chooses \((A_i)_R = (A_i)_G\), then the rotational non-invariance caused by turning on vector fields in the background may be compensated by the global gauge transformation, and the physical gauge configuration may preserve the rotational symmetry [7].

Bearing in mind the above discussion, we work in the temporal gauge \( A_0 = 0 \), and in order to respect for the cosmological principle at the background level, we just allow for \( t \) dependent field configurations

\[
A^a_\mu = \begin{cases} 
\phi(t)\delta^a_i, & \mu = i \\
0 & \mu = 0
\end{cases}. \tag{2.6}
\]

By identifying the gauge indices with the spatial indices, the rotation group \( SU(2) \) is identified with the global part of the gauge group, as discussed above. Note that \( \phi(t) \) by itself is not a scalar under the general coordinate transformations, while

\[
\psi(t) = \frac{\phi(t)}{a(t)}, \tag{2.7}
\]

is indeed a scalar. Using (2.1), the components of the field strengths in this ansatz are obtained

\[
F^a_{0i} = \dot{\phi}\delta^a_i, \quad F^a_{ij} = -g\phi^2\epsilon^a_{ij}. \tag{2.8}
\]

By fixing the gauge as \( A^a_0 = 0 \), the system is left with nine degrees of freedom, \( A^a_i \). The trivial choice for the action is the Yang-Mills action minimally coupled to Einstein gravity. It is known that in any cosmological model subject to a matter field, the necessary condition for a successful inflation is \( \rho + 3p < 0 \) where \( \rho \) and \( p \) are the energy density and pressure of the matter field, respectively. It is easy to see that the Yang-Mills action minimally
coupled to Einstein gravity will not lead to the inflation. This is because as a result of scale invariance of the Yang-Mills action one immediately obtains \( p = \rho/3 \) and \( \rho \geq 0 \), which obviously do not satisfy the above necessary condition. Hence, we need to modify the Yang-Mills action minimally coupled to Einstein gravity. One such appropriate choice has recently been considered in [7].

\[
S = \int d^4x \sqrt{-g} \left[ -\frac{R}{2} - \frac{1}{4} F^a_{\mu\nu} F^a_{\mu\nu} + \frac{\kappa}{384} (\kappa_{\mu\nu})^2 \right], \quad (2.9)
\]

which has led to a successful inflation. The appearance of \( F^4 \) term in this action has been justified by the requirement that the contribution of this term to the energy momentum tensor results in the equation of state \( p = -\rho \), which makes it suitable for driving inflationary dynamics. However, a reasonable justification for this term needs rigorous quantum gauge field theory analysis and some particle physics settings [7].

3 Gauge Inflation by non-Abelian \( F^2 \) term non-minimally coupled to gravity

In this section, we introduce a new action where the Yang-Mills field is non-minimally coupled to Einstein gravity as follows

\[
S = \int d^4x \sqrt{-g} \left[ -\frac{R}{2} - \frac{1}{4} (\kappa g^{\rho\mu} + \alpha G^{\rho\mu}) (\kappa g^{\nu\lambda} + \alpha G^{\nu\lambda}) F_{\rho\lambda} F^a_{\mu\nu} \right], \quad (3.1)
\]

where \( g^{\mu\nu} \) is the metric tensor, \( G^{\mu\nu} \) is the Einstein tensor and \( \kappa > 0, \alpha < 0 \) are constant parameters. In this action, we just use one \( F^2 \) term instead of \( F^2 \) and \( F^4 \) terms considered in the previous action [7]. The price we pay for this replacement is to introduce the Einstein tensor besides the metric tensor as the geometric coupling of two \( F \) terms, namely we work with \( F^2 \) term non-minimally coupled to gravity. We aim to show that it is possible to have inflation in this Yang-Mills field non-minimally coupled to gravity. In this direction, we first obtain the energy momentum tensor

\[
T_{\alpha\beta} = \{ F_{a\alpha} F^a_{\mu\beta} - \frac{1}{4} g_{\alpha\beta} F_{a\mu\nu} F^a_{\mu\nu} \}
\]

\[
+ \frac{\alpha \kappa}{4} \{ g_{\alpha\beta} g^{\rho\mu} G^{\nu\lambda} F_{a\rho\lambda} F^a_{\mu\nu} - 2 G^{\rho\mu} F_{a\rho\beta} F^a_{\mu\alpha} + g_{\alpha\beta} R + \tilde{g}_{\lambda\nu} \tilde{g}^{\lambda\nu} R_{\alpha\beta} \}
\]

\[
+ \Box F(\phi) g_{\alpha\beta} - \nabla_\alpha \nabla_\beta F(\phi) - 4 \tilde{g}_{\lambda\beta} \tilde{R}^\lambda_\alpha \}, \quad (3.2)
\]

where \( F(\phi) = -6(\dot{\phi}^2 - \phi^2 \ddot{\phi}) \) and \( \tilde{g}_{\lambda\nu} = g^{\rho\mu} F_{a\rho\lambda} F^a_{\mu\nu} \). We have ignored the terms including \( \alpha^2 \) in comparison to the terms including \( \alpha \) by assuming \( \alpha \) to be small. This assumption is based on the hypothesis that a considerable contribution of Einstein tensor in the kinetic term of the gauge fields must be limited to the early universe in order to set up an inflationary era. In other words, we expect that the presence of \( \alpha G^{\rho\mu} \) in the kinetic term of the gauge fields is considerable just in the early universe with very small value of the scale factor and large value of \( H \) to trigger the inflation, and it should fade away when the scale factor becomes
large and $H$ becomes small. The presence of small $\alpha$ in front of the Einstein tensor can help to achieve this goal, such that after the inflation is ended the term $\alpha G_{\mu\nu}$ is almost vanishing.

Using the Friedmann-Robertson-Walker (FRW) background (2.2) and the gauge field ansatz (2.6), we may cast the above energy-momentum tensor in the form of a homogenous perfect fluid

$$T^\mu_\nu = \text{diag}(\rho, p, p, p),$$

where

$$\rho = \frac{3\kappa}{2}(\frac{\dot{\phi}^2}{a^2} + \frac{g^2\phi^4}{a^4}) + \frac{\alpha\kappa}{4}\{-3(\frac{\dot{\phi}^2}{a^2} + 2\frac{g^2\phi^4}{a^4})(\frac{2\dot{H} + 3H^2}{a^2}) + 9H^2(\frac{\dot{\phi}^2}{a^2}) - 3(\frac{\ddot{\phi}^2}{a^2} - 2\frac{g^2\phi^4}{a^4})R + 6[\frac{\dot{\phi}^2}{a^2} - \frac{g^2\phi^4}{a^4}]R_0\},$$

(3.3)

and

$$p = \frac{\kappa}{2}(\frac{\dot{\phi}^2}{a^2} + \frac{g^2\phi^4}{a^4}) + \frac{\alpha\kappa}{4}\{3(\frac{\dot{\phi}^2}{a^2} + 2\frac{g^2\phi^4}{a^4})(\frac{2\dot{H} + 3H^2}{a^2}) - 9H^2(\frac{\dot{\phi}^2}{a^2}) - 3(\frac{\ddot{\phi}^2}{a^2} - 2\frac{g^2\phi^4}{a^4})R + 6[\frac{\dot{\phi}^2}{a^2} - \frac{g^2\phi^4}{a^4}]R_i - \ddot{\tilde{F}}(\phi)\},$$

(3.4)

where

$$R = 6[\dot{H} + 2H^2], \quad R_0 = 3[\dot{H} + H^2], \quad R_i = [\dot{H} + 3H^2].$$

(3.5)

We may divide $\rho$ and $p$ in the following way

$$\rho = \rho_{YM} + \rho_\alpha, \quad p = p_{YM} + p_\alpha,$$

(3.6)

where

$$\rho_{YM} = \frac{3\kappa}{2}(\frac{\dot{\phi}^2}{a^2} + \frac{g^2\phi^4}{a^4}), \quad p_{YM} = \frac{\kappa}{2}(\frac{\dot{\phi}^2}{a^2} + \frac{g^2\phi^4}{a^4}),$$

(3.7)

$$\rho_\alpha = \frac{\alpha\kappa}{4}\{-3(\frac{\dot{\phi}^2}{a^2} + 2\frac{g^2\phi^4}{a^4})(\frac{2\dot{H} + 3H^2}{a^2}) + 9H^2(\frac{\dot{\phi}^2}{a^2}) - 3(\frac{\ddot{\phi}^2}{a^2} - 2\frac{g^2\phi^4}{a^4})R + 6[\frac{\dot{\phi}^2}{a^2} - \frac{g^2\phi^4}{a^4}]R_0\},$$

(3.8)

$$p_\alpha = \frac{\alpha\kappa}{4}\{3(\frac{\dot{\phi}^2}{a^2} + 2\frac{g^2\phi^4}{a^4})(\frac{2\dot{H} + 3H^2}{a^2}) - 9H^2(\frac{\dot{\phi}^2}{a^2}) - 3(\frac{\ddot{\phi}^2}{a^2} - 2\frac{g^2\phi^4}{a^4})R + 6[\frac{\dot{\phi}^2}{a^2} - \frac{g^2\phi^4}{a^4}]R_i - \ddot{\tilde{F}}(\phi)\}. $$

(3.9)

Now, we assume that the effective pressure corresponding to the Einstein tensor in the kinetic term is negative as $p_\alpha = -\rho_\alpha$. This gives us the following equation

$$\frac{d^2 x}{dt^2} - \frac{d^2 y}{dt^2} - \frac{1}{3}(\dot{H} + 3H^2)x = 0,$$

(3.10)

where

$$x = \frac{\dot{\phi}^2}{a^2}, \quad y = \frac{g^2\phi^4}{a^4}.$$  

(3.11)

Using the condition $p_\alpha = -\rho_\alpha$ in the corresponding Friedmann equations and also (3.5) we obtain

$$3H^2 = \left(\frac{3\kappa}{2} - \frac{\alpha\kappa}{4}A\right)x + \left(\frac{3\kappa}{2} - \frac{\alpha\kappa}{4}B\right)y,$$

(3.12)
\[ \dot{H} = -\kappa (x + y), \]  
\[ \frac{d}{dt} \left( \frac{3\kappa}{2} (x + y) - \frac{\alpha \kappa}{4} (Ax + By) \right) + 6\kappa H (x + y) = 0, \]  
\[ \text{where} \]
\[ A = 3 \left( \frac{2\dot{H} + 3H^2}{a^2} + 3H^2 \right), \]  
\[ B = 6 \left( \frac{2\dot{H} + 3H^2}{a^2} + R_0^0 \right). \]  
Now, in order to work out with the slow-roll approximation we define the slow-roll parameters
\[ \varepsilon \equiv -\frac{\dot{H}}{H^2} = \frac{3}{2} \frac{\rho + p}{\rho}, \quad \eta = \varepsilon - \frac{\dot{\varepsilon}}{2H\varepsilon}, \]  
from which one may specify the slow-roll dynamics by \( \varepsilon, \eta \ll 1 \). Using the Friedmann equations
\[ 3H^2 = \rho = (\rho_{YM} + \rho_\alpha), \]  
\[ 2\dot{H} + 3H^2 = -p = -(p_{YM} + p_\alpha), \]  
where \( p_\alpha = -\rho_\alpha \) and \( p_{YM} = \frac{1}{3}\rho_{YM} \), we obtain
\[ \varepsilon = \frac{2\rho_{YM}}{\rho_{YM} + \rho_\alpha}. \]  
Therefore, to have the slow-roll approximation the \( \alpha \)-term contribution \( \rho_\alpha \) should dominate over the Yang-Mills contribution \( \rho_{YM} \), namely \( \rho_\alpha \gg \rho_{YM} \). This slow-roll approximation is compatible with the condition \( p_\alpha = -\rho_\alpha \). In fact, considering Eq.(3.8) and Eq.(3.9), we realize that although \( \alpha \) is small, but at the beginning of inflation \( H \) is so large and the scale factor \( a \) is so small that the first term of \( \rho_\alpha \) in (3.8) containing \( \frac{2H + 3H^2}{a^2} \) can dominate over the other terms, and even dominates \( \rho_{YM} \) defined in Eq.(3.7). Hence, during the inflation it is reasonable to consider \( \rho_\alpha \gg \rho_{YM} \), and \( |p_\alpha| \gg |\rho_{YM}| \). This confirms the compatibility of the condition \( p_\alpha = -\rho_\alpha \) with the slow-roll approximation \( \rho_\alpha \gg \rho_{YM} \).

However, note that it is not enough to make sure \( \varepsilon \ll 1 \) for a successful inflation. In fact, time-variations of \( \varepsilon \) and all the other physical dynamical variables of the problem, like \( \eta \) and the \( \psi \) field, must also remain small during the inflation over a reasonable period of time. For typical slow-roll models, \( \eta \) usually measures the rolling velocity of the inflaton, namely \( \eta = -\frac{\dot{\psi}}{H\psi} \). Hence, it is useful to define
\[ \delta \equiv -\frac{\dot{\psi}}{H\psi}, \]
which is related to $\varepsilon$ and $\eta$ through the equations

\begin{align}
\varepsilon &= 2 + \frac{\alpha \kappa}{6} \left\{(1 - \delta)^2 A + B \gamma \right\} \psi^2, \\
\eta &= \varepsilon - (2 - \varepsilon) \left(\frac{\delta}{\varepsilon} - \frac{\dot{Z}}{2HZ\varepsilon}\right),
\end{align}

and

\begin{align}
Z &= A + \gamma B, \\
\gamma &= \frac{g^2 \psi^2}{H^2}.
\end{align}

The slow-roll regime requires $\dot{\varepsilon} \sim H\varepsilon^2$ and $\eta \sim \varepsilon$, so we demand in (3.20b) that

\begin{align}
\delta &\sim \varepsilon^2, \\
\frac{\dot{Z}}{2HZ} &\sim \varepsilon^2.
\end{align}

Now, using $\delta \sim \varepsilon^2 \approx 0$ and (3.21a) in (3.20a) we have

\begin{align}
\varepsilon &= 2 + \frac{\alpha \kappa}{6} Z \psi^2.
\end{align}

Using Eqs. (3.11), (3.12), (3.17), (3.19), (3.21b) and $\delta \approx 0$ we obtain

\begin{align}
\rho &= \left(\frac{3\kappa}{2} - \frac{\alpha \kappa}{4} A\right) H^2 \psi^2 + \left(\frac{3\kappa}{2} - \frac{\alpha \kappa}{4} B\right) \gamma H^2 \psi^2,
\end{align}

or

\begin{align}
\rho &\approx H^2 \psi^2 \left[\frac{3\kappa}{2} (1 + \gamma) - \frac{\alpha \kappa}{4} Z\right].
\end{align}

Explicitly, the equations (3.17), (3.20b), (3.23), and (3.25) admit the solutions

\begin{align}
\varepsilon &\approx \kappa \psi^2 (\gamma + 1), \\
\delta &= \frac{\dot{x}}{2}(\varepsilon - \eta) + \frac{\dot{Z}}{2HZ},
\end{align}

where $\approx$ means equality to first order in the slow-roll parameter $\varepsilon$ and

\begin{align}
\gamma &= \frac{g^2 \psi^2}{H^2}, \quad \iff \quad H^2 \approx \frac{g^2 \varepsilon}{\kappa \gamma (\gamma + 1)}.
\end{align}

Using $p_\alpha = -\rho_\alpha$, the conservation equation for the perfect fluid becomes

\begin{align}
\dot{\rho} + 3H(\rho_{YM} + p_{YM}) = 0.
\end{align}

Note that, $p_\alpha = -\rho_\alpha$ means that (3.10) and (3.11) are compatible with the equation of motion for the gauge field (i.e. $\phi$ field). Putting (3.25) in (3.28) and using (3.7) and $\varepsilon^2 \approx 0$ yields

\begin{align}
\dot{\gamma} - \frac{4H}{9\kappa} (1 + \gamma) &\approx 0,
\end{align}
which, using Eqs. (3.21b) and (3.26a), is rewritten as
\begin{equation}
\dot{\gamma} - \frac{4g}{9\kappa} \sqrt{\frac{\varepsilon(1 + \gamma)}{\kappa\gamma}} \simeq 0.
\end{equation}

Because $g$ and $\varepsilon$ are small, we have
\begin{equation}
\dot{\gamma} \approx 0,
\end{equation}
which shows that $\gamma$ is a positive parameter which is slowly varying during the slow-roll regime. Considering (3.21b) with $H \approx \text{const}$ and that $\gamma$ is a slowly varying parameter, we find that $\psi$ is slowly varying, too. Hence, (3.21b) implies that $\gamma H^2$ is almost constant during the slow-roll regime from which we obtain
\begin{equation}
\frac{\gamma}{\gamma_i} \simeq \frac{H_f^2}{H_i^2} \iff \frac{\varepsilon}{\varepsilon_i} \simeq \frac{\gamma + 1}{\gamma_i + 1},
\end{equation}
where $\varepsilon_i$, $\gamma_i$ and $H_i$ are the values of parameters at the beginning of inflation. It is obvious that the slow-roll regime ends when $\varepsilon_f \approx 1$. This happens when the scale factor becomes inflationary large and $H$ becomes small towards the end of inflation and the contribution of $\alpha G^{\mu\nu}$ in the kinetic term of the gauge fields starts decreasing such that $\rho_\alpha \sim \rho_{YM}$. Therefore, we find
\begin{equation}
\gamma_f \simeq \frac{\gamma_i + 1}{\varepsilon_i}, \quad \frac{H_f^2}{H_i^2} \simeq \frac{\gamma_i}{\gamma_i + 1} \varepsilon_i.
\end{equation}

Now, we can compute the number of e-folding
\begin{equation}
N_e = \int_{t_i}^{t_f} H dt = - \int_{H_i}^{H_f} \frac{dH}{\varepsilon H} \simeq \frac{\gamma_i + 1}{2\varepsilon_i} \ln \frac{\gamma_i + 1}{\gamma_i},
\end{equation}
where use has been made of $\varepsilon \equiv -\frac{H}{H_i^2}$. It is seen that number of e-folding is inversely related to the slow roll parameter $\varepsilon_i$, so the more small slow roll parameter, the more e-folding.

### 4 Conclusion

In this paper, we have combined two recently developed ideas of *Gauge-flation* and *Kinetic coupled gravity*. Gauge-flation has been introduced to establish inflation in the context of high energy physics where the slow-roll regime is set up by a non-Abelian gauge field $A$ with the field strength $F$. This class of models are based on a gauge field theory with $F^2$ and $F^4$ terms with a non-Abelian gauge group minimally coupled to gravity. Kinetic coupled gravity, on the other hand, has been introduced mainly to account for the current acceleration of the universe. The non-minimal kinetic coupled gravity is set up by introducing the Einstein tensor besides the metric tensor within the $F^2$ term. By combining these ideas, we have presented a new class of inflationary models based on a non-Abelian gauge field theory with $F^2$ term non-minimally coupled to gravity.
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