

D-term chaotic inflation in supergravity

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Even though the chaotic inflation is one of the most popular inflation models for its simple dynamics and compelling resolutions to the initial condition problems, its realization in supergravity has been considered a challenging task. We discuss how the chaotic inflation dominated by the D-term can be induced in supergravity, which would give a new perspective on the inflation model building in supergravity.

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Cosmic inflation has been one of the most successful early universe scenarios for more than 25 years, with the ever-growing supports from the observations of the cosmic microwave background anisotropies and the large scale structure of the universe [1, 2, 3]. Among many types of inflation models proposed so far, chaotic inflation is special for its amelioration of the initial condition problems [4] and it would be of considerable interest to realize chaotic inflation in a sensible particle physics theory.

One of the leading theories as an extension of the minimal standard model is supersymmetry [5] which gives the attractive solutions to the hierarchy problem of the standard model as well as the unification of three gauge couplings. In particular, in the early universe, its local version supergravity would govern the dynamics of the universe, while there has been criticisms for the implementation of chaotic inflation in supergravity. This is simply because the scalar potential coming from the F-term has an exponential dependence on the Kähler potential. This prevents the scalar fields from acquiring the amplitudes larger than the reduced Planck scale $M_p \simeq 2.4 \times 10^{18} \text{GeV}$ and also spoils the flatness of an inflaton potential (so called η -problem). In order to circumvent this difficulty [6], Kawasaki, Yanagida and one of the present authors (M.Y.) introduced the Nambu-Goldstone-like shift symmetry [7], so that the imaginary part of the an inflaton field does not suffer from the exponential growth. Though such a shift symmetry is motivated from the string theory and is similar to, for example, the Heisenberg symmetry [8], it may be difficult to associate it with the low energy effective theory of particle physics such as grand unified theory (GUT). It would be then worth seeking an alternative supergravity chaotic inflation model without such a symmetry possibly absent in the effective field theory.

The scalar potential in supergravity also consists of, in addition to the F-term, the D-term which does not have an aforementioned dangerous exponential factor. A plausible possibility to realize the supergravity chaotic infla-

tion then would be to consider the inflation models where the D-term dominates over the F-term. In the conventional D-term inflation models [9, 10, 11, 12], the energy density is sourced by the Fayet-Iliopoulos (FI) term ξ in D-term and the slope of an inflaton potential is induced by the one-loop corrections. Since the one-loop corrections cannot exceed the tree level potential energy density of order $\xi^2 \ll M_p^4$, the inflation cannot start from the Planckian energy scale. This in turn implies that chaotic inflation is not compatible with the standard D-term inflation models because the universe would collapse before such an inflation energy scale is reached unless the universe started with an open geometry. It would be then intriguing to consistently incorporate chaotic inflation in a D-term dominated inflation model, possibly without making use of the FI term dominance or one-loop corrections.

In this paper, we present a chaotic inflation model in supergravity where the D-term dominates the inflation energy density. The D-term is helpful for the realization of chaotic inflation by allowing the (beyond) Planckian amplitude for an inflaton field as well as for the avoidance of the η -problem. Our D-term chaotic inflation model requires neither the shift symmetry nor the dominance of FI term. No need for the FI term dominance in our model alleviates the potentially unnatural tuning of its amplitude ξ which otherwise would be constrained tightly by the cosmic perturbations and cosmic strings as in the conventional D-term inflation models [13].

We introduce four superfields Φ_i ($i = 1, 2, 3, 4$) charged under $U(1)$ gauge symmetry and (global) $U(1)_R$ symmetry. The charges Q_i, Q_i^R of the superfields are given in Table I which ensure our toy model is anomaly free [14, 15]. Then, the general (renormalizable) superpotential for these fields is given by

$$W = a\Phi_1^2\Phi_3 - b\Phi_2\Phi_3 + c\Phi_2\Phi_4^2, \quad (1)$$

where we set the constants a, b, c to be real and positive for simplicity and a non-renormalizable term $\Phi_1^2\Phi_4^2$ is omitted since it does not change the dynam-

	Φ_1	Φ_2	Φ_3	Φ_4
Q	+1	+2	-2	-1
Q^R	0	0	+2	+1

TABLE I: The charges of various superfields of $U(1) \times U(1)_R$.

ics essentially. Taking the canonical Kähler potential, $K(\Phi_i, \Phi_i^*) = \sum_i |\Phi_i|^2$, and the minimal gauge kinetic function, $f_{ab}(\Phi_i) = \delta_{ab}$, leads to the scalar potential consisting of the F-term V_F and D-term V_D

$$\begin{aligned}
V &= V_F + V_D, \\
V_F &= e^K \left[|2a\phi_1\phi_3 + \phi_1^*W|^2 + |-b\phi_3 + c\phi_4^2 + \phi_2^*W|^2 \right. \\
&\quad \left. + |a\phi_1 - b\phi_2 + \phi_3^*W|^2 + |2c\phi_2\phi_4 + \phi_4^*W|^2 - 3|W|^2 \right], \\
V_D &= \frac{g^2}{2} (|\phi_1|^2 + 2|\phi_2|^2 - 2|\phi_3|^2 - |\phi_4|^2)^2, \quad (2)
\end{aligned}$$

where g is the gauge coupling of the $U(1)$ symmetry, and we take the vanishing FI term for simplicity. Here and hereafter we set the reduced Planck scale M_p to be unity.¹

The minimum of the F-term (the F-flat condition) is given by

$$V_F = 0 \iff \begin{cases} a\phi_1^2 - b\phi_2 = 0 & \& \phi_3 = \phi_4 = 0. \\ & \text{or} \\ \phi_1 = \phi_2 = 0 & \& -b\phi_3 + c\phi_4^2 = 0. \end{cases} \quad (3)$$

On the other hand, the minimum of the D-term (the D-flat condition) is given by

$$V_D = 0 \iff |\phi_1|^2 + 2|\phi_2|^2 = 2|\phi_3|^2 + |\phi_4|^2. \quad (4)$$

Hence, the global minimum of the potential is given by

$$\phi_1 = \phi_2 = \phi_3 = \phi_4 = 0. \quad (5)$$

However, when the universe starts around the Planck scale, if $|\phi_1| \gg 1$ or $|\phi_2| \gg 1$, and $|\phi_3|, |\phi_4| \lesssim 1$ for example, the almost F-flat condition ($a\phi_1^2 \simeq b\phi_2, \phi_3 = \phi_4 = 0$) is first realized due to the exponential factor e^K of the F-term. Consequently, the potential is mostly dominated by the D-term and chaotic inflation can take place.

Because the system is invariant under the following transformation

$$\phi_1 \leftrightarrow \phi_4, \quad \phi_2 \leftrightarrow \phi_3, \quad a \leftrightarrow c, \quad Q \rightarrow -Q, \quad (6)$$

¹ The qualitative features of our chaotic inflation model would not be affected even if the non-renormalizable terms, such as $\lambda\phi^4(\phi/M_p)^n$, appear in the potential as long as $\lambda \ll 1$, that is, the effective cut-off scale is larger than the reduced Planck scale. Such a small λ associated with the breaking of the R-symmetry used in our toy model to prohibit the non-renormalizable terms is natural in 't Hooft's sense.

the dynamics is essentially the same even if we interchange ϕ_1 and ϕ_2 by ϕ_3 and ϕ_4 . Thus, we concentrate on the case that $|\phi_1| \gg 1$ or $|\phi_2| \gg 1$, and $|\phi_3|, |\phi_4| \lesssim 1$.

Now, we investigate the dynamics in details. Despite the e^K factor of F-terms, due to the presence of the relatively rather small but non-vanishing D-terms, the actual inflaton trajectory is slightly deviated from the exact F-flat direction and given by solving the equations (1) $\partial V/\partial\phi_i^* = 0$ ($i = 2, 3, 4$) or (2) $\partial V/\partial\phi_i^* = 0$ ($i = 1, 3, 4$), depending on the magnitude of b/a , as clarified later. The solution (named M1) of the first equations is given by $\phi_2 = \phi_2(\phi_1), \phi_3 = \phi_4 = 0$ and that (named M2) of the second equations is given by $\phi_1 = \phi_1(\phi_2), \phi_3 = \phi_4 = 0$.

Here, we check whether inflation indeed occurs along this field trajectory. For this purpose, we first evaluate the mass terms of the fields ϕ_3 and ϕ_4 along these trajectories M1 and M2, $V_{,ij}|_M \phi_i^* \phi_j$ with $V_{,ij} \equiv \partial^2 V / (\partial\phi_i^* \partial\phi_j)$. The suffix M represents the evaluations along either the trajectory M1 or M2. Then, the mass matrix of the fields ϕ_3 and ϕ_4 , $V_{,ij}|_M$, is given by

$$\begin{aligned}
V_{,33}|_M &\simeq e^K (4a^2|\phi_1|^2 + b^2), \quad V_{,3i}|_M = 0 \quad (\text{for } i = 1, 2, 4), \\
V_{,44}|_M &\simeq 4e^K c^2|\phi_2|^2, \quad V_{,4i}|_M = 0 \quad (\text{for } i = 1, 2, 3). \quad (7)
\end{aligned}$$

Both of these masses are much larger than $H^2 \simeq g^2(|\phi_1|^2 + 2|\phi_2|^2)^2/2$ unless the constants a, b, c are exponentially small (typically $e^{-g^{-1}}, g = \mathcal{O}(10^{-6})$), which makes ϕ_3 and ϕ_4 quickly go to the zeros. As a result, we can safely set ϕ_3 and ϕ_4 to be zero and we can discuss the dynamics of the inflaton based on the following potential,

$$V_{\text{eff}}(\phi_1, \phi_2) \equiv V(\phi_1, \phi_2, \phi_3 = 0, \phi_4 = 0). \quad (8)$$

By use of the $U(1)$ gauge symmetry, we can, for instance, make the field ϕ_1 real without loss of generality, so that the $\text{Im}\phi_2$ rapidly goes to the zero because the effective mass squared of the imaginary part of ϕ_2 is given by $m_{\text{Im}\phi_2}^2 \simeq b^2 e^K$. We therefore consider the following effective potential by redefining the fields $\phi_i \equiv \sqrt{2} \text{Re} \phi_i$ ($i = 1, 2$, and we take ϕ_i to be positive for definiteness) and $b' \equiv \sqrt{2}b$,

$$V_{\text{eff}}(\phi_1, \phi_2) = \frac{1}{4} e^K (a\phi_1^2 - b'\phi_2)^2 + \frac{g^2}{8} (\phi_1^2 + 2\phi_2^2)^2 \quad (9)$$

with $K = (\phi_1^2 + \phi_2^2)/2$ and the canonical kinetic terms.

Now, we would like to discuss the dynamics of the inflation based on the above potential. First of all, we consider the region where $B \equiv b'/a \gg \phi_1 (\gg 1)$ is satisfied. As shown later, in this region, the field ϕ_1 plays the role of the inflaton while the inflationary trajectory is almost determined by the condition (M1) $\partial V_{\text{eff}}/\partial\phi_2 = 0$, which is equivalent to

$$\begin{aligned}
a\phi_1^2 - b'\phi_2 &= e^{-K} \frac{g^2(\phi_1^2 + 2\phi_2^2)\phi_2}{\frac{1}{2}b' - \frac{1}{4}\phi_2(a\phi_1^2 - b'\phi_2)} \\
&\simeq e^{-K} \frac{2}{b'} g^2(\phi_1^2 + 2\phi_2^2)\phi_2. \quad (10)
\end{aligned}$$

Here note that combining $a\phi_1^2 - b'\phi_2 = \mathcal{O}(e^{-K})$ with $B = b'/a \gg \phi_1$ leads to $\phi_1 \gg \phi_2$. The F-term contribution to the potential is estimated as

$$V_F = \frac{\frac{1}{4}\phi_2(a\phi_1^2 - b'\phi_2)}{\frac{1}{2}b' - \frac{1}{4}\phi_2(a\phi_1^2 - b'\phi_2)} g^2(\phi_1^2 + 2\phi_2^2) < g^2(\phi_1^2 + 2\phi_2^2) \ll V_D = \frac{g^2}{2}(\phi_1^2 + 2\phi_2^2)^2 \quad (11)$$

for $\phi_1 \gg 1$. Thus, the potential is dominated by the D-term during inflation. We also consider the mass matrix of the fields ϕ_1 and ϕ_2 , $V_{,ij}|_{M1} (V_{,ij} \equiv \partial^2 V_{\text{eff}}/(\partial\phi_i\partial\phi_j))$, which is given by

$$\begin{aligned} V_{,11}|_{M1} &\simeq 2a^2\phi_1^2 e^K + 2\frac{a}{b'}g^2(\phi_1^2 + 2\phi_2^2)\phi_2(1 + 2\phi_1^2) \\ &\quad + \frac{g^2}{2}(3\phi_1^2 + 2\phi_2^2), \\ V_{,12}|_{M1} &\simeq -ab'\phi_1 e^K + 2\frac{a}{b'}g^2(\phi_1^2 + 2\phi_2^2)\phi_1\phi_2(\phi_2 - \frac{b'}{2a}) \\ &\quad + 2g^2\phi_1\phi_2, \\ V_{,22}|_{M1} &\simeq \frac{1}{2}b'^2 e^K - 2g^2(\phi_1^2 + 2\phi_2^2)\phi_2^2 + g^2(\phi_1^2 + 6\phi_2^2)12 \end{aligned}$$

up to the order of $\mathcal{O}(e^K)^0$.² The effective mass squared λ of the fields ϕ_1 and ϕ_2 is given as the solutions of the following equation,

$$\lambda^2 - (V_{,11} + V_{,22})\lambda + V_{,11}V_{,22} - V_{,12}^2 = 0, \quad (13)$$

where

$$\begin{aligned} V_{,11} + V_{,22}|_{M1} &\simeq e^K \left(2a^2\phi_1^2 + \frac{1}{2}b'^2 \right) \simeq \frac{1}{2}b'^2 e^K, \\ V_{,11}V_{,22} - V_{,12}^2|_{M1} &\simeq \frac{3}{4}g^2b'^2\phi_1^2 e^K \end{aligned} \quad (14)$$

up to the order of $\mathcal{O}(e^K)$. Here, we have used the approximation that $a\phi_1^2 - b'\phi_2 = \mathcal{O}(e^{-K})$, $B = b'/a \gg \phi_1$, and $\phi_1 \gg \phi_2$. The effective squared masses are then approximately given by

$$\lambda \simeq \frac{1}{2}b'^2 e^K, \quad \frac{3}{2}g^2\phi_1^2 \ll H^2 \simeq V_D/3, \quad (15)$$

where H is a Hubble parameter. The inflaton field in the chaotic inflation corresponds to this effectively massless mode. This light mass squared vanishes for $g = 0$ as expected, reflecting the exact F-flat direction.

Since $V_{,22} \gg V_{,11}$ and $\phi_1 \gg \phi_2$, the inflationary trajectory is given by the minimum of the field ϕ_2 , $\partial V/\partial\phi_2 = 0$, which enables us to write the minimum of ϕ_2 as a function of ϕ_1 , $\phi_2^m = \phi_2^m(\phi_1)$. The field trajectory governing the inflation dynamics therefore can be parameterized by

the field ϕ_1 which we call an inflaton.³ Then, by inserting the above relation into the effective potential, we define the reduced potential $V_{r1}(\phi_1)$ as

$$V_{r1}(\phi_1) \equiv V_{\text{eff}}(\phi_1, \phi_2^m(\phi_1)) \left(\simeq \frac{g^2}{8}\phi_1^4 \right). \quad (16)$$

As explicitly shown in Ref. [16], when there is only one massless mode and the other modes are massive, the generation of adiabatic density fluctuations as well as the dynamics of the homogeneous mode is completely determined by the reduced potential $V_{r1}(\phi_1)$. Indeed, the equation of motion for the homogeneous mode of the inflaton ϕ_1 along the rolling direction (M1) is approximated as

$$\ddot{\phi}_1 + 3H\dot{\phi}_1 + \left. \frac{\partial V_{\text{eff}}}{\partial\phi_1} \right|_{M1} = \ddot{\phi}_1 + 3H\dot{\phi}_1 + \frac{dV_{r1}}{d\phi_1} = 0, \quad (17)$$

where the dot represents time derivative. Thus, the dynamics of the inflation with the inflaton ϕ_1 can be estimated by using the reduced potential $V_{r1}(\phi_1)$ as long as the dynamics rolls down along the minimum of ϕ_2 (M1).

Next, we evaluate the primordial density fluctuations in the longitudinal gauge. The equation of motion for the perturbation $\delta\phi_i$ of each real field is given by [17]

$$\begin{aligned} \ddot{\delta\phi}_i + 3H\dot{\delta\phi}_i - \frac{\nabla^2}{a^2}\delta\phi_i + \sum_j \left. \frac{\partial^2 V_{\text{eff}}}{\partial\phi_j\partial\phi_i} \right|_{M1} \delta\phi_j \\ = -2 \left. \frac{\partial V_{\text{eff}}}{\partial\phi_i} \right|_{M1} \Phi + 4\dot{\phi}_i\dot{\Phi}, \end{aligned} \quad (18)$$

where Φ is the gravitational potential. We hereafter use the same symbol ϕ_i for both the homogeneous mode and the full field for notational brevity unless stated otherwise.

We are interested only in the adiabatic density fluctuations characterized by the condition

$$\frac{\dot{\delta\phi}_1}{\dot{\phi}_1} = \frac{\dot{\delta\phi}_2}{\dot{\phi}_2} \iff \delta\phi_2 = \frac{d\phi_2^m(\phi_1)}{d\phi_1}\delta\phi_1 \quad (19)$$

where we have used $\dot{\phi}_2/\dot{\phi}_1 = d\phi_2^m(\phi_1)/d\phi_1$. Since the relation $\partial V_{\text{eff}}(\phi_1, \phi_2^m(\phi_1))/\partial\phi_2 = 0$ holds for any ϕ_1 in the relevant region, we find

$$\frac{d}{d\phi_1} \left[\left. \frac{\partial V_{\text{eff}}}{\partial\phi_2}(\phi_1, \phi_2^m(\phi_1)) \right] = V_{,12} + V_{,22} \frac{d\phi_2^m}{d\phi_1} \Big|_{M1} = 0. \quad (20)$$

Taking into account this relation and

$$\frac{d^2 V_{r1}}{d\phi_1^2} = V_{,11} + 2 \frac{d\phi_2^m}{d\phi_1} V_{,12} + \left(\frac{d\phi_2^m}{d\phi_1} \right)^2 V_{,22} \Big|_{M1}$$

² More precisely, we have used the approximation that $b' \gg |\phi_2(a\phi_1^2 - b'\phi_2)|$ and $a, b'\phi_2 \gg |a\phi_1^2 - b'\phi_2|$.

³ Note here that the effectively massless field trajectory parameterized by the inflaton field ϕ_1 is different from the ϕ_1 direction with the mass $V_{,11} \gg H^2$.

$$= \frac{V_{,11}V_{,22} - V_{,12}^2}{V_{,22}} \Big|_{M_1} \left(\simeq \frac{3}{2}g^2\phi_1^2 \right), \quad (21)$$

the equation of motion for the perturbation $\delta\phi_1$ becomes

$$\ddot{\delta\phi}_1 + 3H\dot{\delta\phi}_1 - \frac{\nabla^2}{a^2}\delta\phi_1 + \frac{d^2V_{r1}}{d\phi_1^2}\delta\phi_1 = -2\frac{dV_{r1}}{d\phi_1}\dot{\Phi} + 4\dot{\phi}_1\dot{\Phi} \quad (22)$$

Thus, the perturbation $\delta\phi_1$ is completely determined by the reduced potential $V_{r1}(\phi_1)$. Note that $d^2V_{r1}/d\phi_1^2$ is the effective mass squared along the rolling direction given by $\partial V_{\text{eff}}/\partial\phi_2 = 0$. Therefore, this rolling direction actually coincides with the eigenvector of the effectively massless mode of λ .

On the other hand, by use of the adiabatic condition, the gravitational potential is described only by $\delta\phi_1$ as

$$\left(\dot{H} - \frac{\nabla^2}{a^2} \right) \Phi = \frac{1}{2M_G^2} \left(\ddot{\phi}_1\delta\phi_1 - \dot{\phi}_1\dot{\delta\phi}_1 \right) \left[1 + \left(\frac{d\phi_2^m}{d\phi_1} \right)^2 \right] \quad (23)$$

In the long wave limit, $\Phi \simeq (d/dt)(\delta\phi_1/\dot{\phi}_1)$, where we have used $\dot{H} = -(\dot{\phi}_1^2 + \dot{\phi}_2^2)/(2M_G^2) = -\dot{\phi}_1^2[1 + (d\phi_2^m/d\phi_1)^2]/(2M_G^2)$. This expression of the gravitational potential coincides with that of the single field inflation with the reduced potential $V_{r1}(\phi_1)$. We can in consequence calculate the density fluctuations of our inflation model based on the reduced potential $V_{r1}(\phi_1) \simeq g^2\phi_1^4/8$.

In the opposite region where $B = b/a \ll \phi_1 (\ll \phi_2)$, the inflationary trajectory is given by the minimum of ϕ_1 , $\phi_1^m = \phi_1^m(\phi_2)$ coming from the condition (M2) $\partial V_{\text{eff}}/\partial\phi_1 = 0$. The field ϕ_2 now plays a role of the inflaton. Defining the reduced potential $V_{r2}(\phi_2) \equiv V_{\text{eff}}(\phi_1^m(\phi_2), \phi_2)$, we can easily show that the generation of adiabatic density fluctuations as well as the dynamics of the homogeneous mode can be calculated by using the reduced potential $V_{r2}(\phi_2)$.

For $B > 1$, in the region where $B \sim \phi_1$, the use of the condition $\partial V_{\text{eff}}/\partial\phi_1 = 0$ or $\partial V_{\text{eff}}/\partial\phi_2 = 0$ may generate the error of the effective mass squared up to the factor 2 due to $V_{,11} \sim V_{,22}$. However, since the true trajectory is effectively massless, we can see from the mass matrix that the deviation is at most $\mathcal{O}(e^{-K})$. It therefore still gives an approximate trajectory of the inflaton, depending on $\phi_1 > \phi_2$ or $\phi_1 < \phi_2$ and hence the qualitative behaviors of the inflaton dynamics and the density fluctuations do not change in this region. Hence the chaotic inflation can be induced in a wide range of the parameters in our toy model. It could also indicate that, once a model possesses a F-flat direction lifted by a D-term, the chaotic inflation could be induced without so much restrictions on the model parameters besides those from the cosmic perturbations.

Finally we give a few comments on our model. The standard procedure shows that the gauge coupling g should be $g \sim 10^{-6}$ in order to explain the primordial

density fluctuations. This value of the gauge coupling is much smaller than the standard gauge couplings. However, this may not be a problem because the gauge symmetry may be a hidden gauge symmetry, or the gauge coupling could be suppressed, for instance, by considering the extra dimensions. Next, even though we presented a toy model of the quartic potential chaotic inflation, the leading order polynomial can be different by an appropriate choice of the non-minimal gauge kinetic function (for instance, a form $f = 1 + d_i|\phi_i|^2$ (d_i : const) could lead to a quadratic potential). Note this is in contrast to the standard D-term inflation with the dominance of the FI term where a non-minimal gauge kinetic function spoils the flatness of the inflaton potential [18]. As for the reheating, it may require the spontaneous breaking of the gauge symmetry after inflation because the inflaton cannot directly decay into the standard particles due to the charge conservation. Such a breaking can occur, for instance, by the introduction of the FI term or a Higgs-like field, and such modifications of our simple toy model presented in this paper would lead to a variety of inflation models with the potentially rich phenomena.⁴ Further study will be given in the forthcoming paper [19].

In summary, we have shown that chaotic inflation can take place in supergravity even without the shift symmetry where the inflaton field trajectory follows the (almost) F-flat direction lifted by the D-term. We stress that, according to our analysis demonstrated through a toy model, many of such F-flat directions could potentially cause the D-term chaotic inflation. The D-term is related to the gauge coupling and a variant of our model could lead to a possible link between a chaotic inflation model and the low energy effective theory of particle physics [20].

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⁴ For instance, an inflaton superpotential of form $W = (a\Phi_1^2 - b\Phi_2^2)\Phi_3$ or $W = M(a\Phi_1 - b\Phi_2)\Phi_3$, with a non-vanishing FI term, could lead to a simpler chaotic inflation dynamics with a subsequent hybrid inflation stage, even though one needs to introduce the additional fields to cancel anomaly.

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