TRAPPED QUINTESSENTIAL INFLATION
FROM FLUX COMPACTIFICATIONS

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Quintessential inflation is studied using a string modulus as the inflaton–
quintessence field. It is assumed that the modulus crosses an enhanced sym-
metry point (ESP) in field space. Particle production at the ESP temporarily
traps the modulus resulting in a period of inflation. After reheating, the modu-
lus freezes due to cosmological friction at a large value, such that its scalar po-
tential is dominated by contributions due to fluxes in the extra dimensions. The
modulus remains frozen until the present, when it can become quintessence.

Keywords: String moduli, inflation, quintessence

1. Introduction

A plethora of observations concur that the Universe at present enters a
phase of accelerated expansion. In fact, most cosmologists accept that over
70% of the Universe content at present corresponds to the elusive dark
energy: a substance with pressure negative enough to cause the observed
acceleration. The simplest form of dark energy is a positive cosmological
constant $\Lambda$, which however, needs to be incredibly fine-tuned to explain
the observations. This is why theorists have looked for alternatives, which
could explain the observations while setting $\Lambda = 0$, as was originally as-
sumed. A promising idea is to consider that the Universe at present is
entering a late-time inflationary period. The credibility of this option is
supported also by the fact that the generic predictions of inflation in the
eyear Universe are in excellent agreement with the observations. The scalar
field responsible for this late-inflation period is called quintessence because
it is the fifth element after baryons, photons, CDM and neutrinos.

Since they are based on the same idea, it is natural to attempt to unify
early Universe inflation with quintessence. Quintessential inflation was thus
This attempt has many advantages. Firstly, quintessential inflation models allow the treatment of both inflation and quintessence within a single theoretical framework. Also, quintessential inflation dispenses with the tuning problem of the initial conditions for quintessence. Finally, unified models for inflation and quintessence are more economic because they avoid introducing yet another unobserved scalar field.

For quintessential inflation to work one needs a scalar field with a runaway potential, such that the minimum has not been reached until today and, therefore, there is residual potential density, which can cause the observed accelerated expansion. String moduli fields are suitable because they are typically characterised by such runaway potentials. The problem with such fields, however, is how to stabilise them temporarily, in order to use them as inflatons in the early Universe. In this work (see also Ref. 9) we achieve this by considering that, during its early evolution our modulus crosses an enhanced symmetry point (ESP) in field space. When this occurs the modulus is trapped temporarily at the ESP,\(^\text{10}\) which leads to a period of inflation. After inflation the modulus picks up speed again in field space resulting into a period of kinetic density domination (kination).\(^\text{11}\) Kination ends when the thermal bath of the hot big bang (HBB) takes over. During the HBB, due to cosmological friction,\(^\text{12}\) the modulus freezes at some large value and remains there until the present, when its potential density dominates and drives the late-time accelerated expansion.\(^\text{8}\)

It is evident that, in order for the modulus to become quintessence, it should not decay after the end of inflation. Reheating, therefore should be achieved by other means. We assume that the thermal bath of the HBB is due to the decay of some curvaton field\(^\text{13}\) as suggested in Refs.\(^\text{8},\text{14}\) By considering a curvaton we do not add an ad hoc degree of freedom, because the curvaton can be a realistic field, already present in simple extensions of the standard model (e.g. a right-handed sneutrino,\(^\text{15}\) a flat direction of the (N)MSSM\(^\text{16}\) or a pseudo Nambu-Goldstone boson\(^\text{17},\text{18}\) possibly associated with the Peccei-Quinn symmetry\(^\text{19}\)). Apart from reheating, the curvaton can provide the correct amplitude of curvature perturbations in the Universe. Consequently, the energy scale of inflation can be much lower than the grand unified scale.\(^\text{20}\) In fact, in certain curvaton models, the Hubble scale during inflation can be as low as the electroweak scale.\(^\text{18},\text{21}\)

2. The runaway scalar potential

String theories contain a number of flat directions which are parametrised by the so-called moduli fields, which correspond to the size and shape of the
compactified extra dimensions. Many such flat directions are lifted by non-perturbative effects, such as gaugino condensation or D-brane instantons.\textsuperscript{22} The superpotential, then, is of the form

\[ W = W_0 + W_{np} \]

where \( W_0 \approx \text{const.} \) is the tree level contribution from fluxes, \( A \) and \( c \) are constants and \( T \) is a Kähler modulus in units of \( m_P \). Hence, the non-perturbative superpotential \( W_{np} \) results in a runaway scalar potential characteristic of string compactifications. For example, in type IIB compactifications with a single Kähler modulus, \( \sigma \equiv \text{Re}(T) \) is the so-called volume modulus, which parametrises the volume of the compactified space. In this case, the runaway behaviour leads to decompactification of the internal manifold. The tree level Kähler potential for a modulus, in units of \( m_P^2 \), is

\[ K = -3 \ln (T + \bar{T}) = -3 \ln(2\sigma), \]

and the corresponding supergravity potential is\textsuperscript{a}

\[ V_{np}(\sigma) \simeq \frac{cAe^{-c\sigma}}{2\sigma^2 m_P^2} \left( \frac{c\sigma}{3} Ae^{-c\sigma} - W_0 \right). \]

To study the cosmology, we turn to the canonically normalised modulus \( \phi \) which, due to Eq. (2), is associated with \( \sigma \) as

\[ \sigma(\phi) = \exp \left( \sqrt{\frac{2}{3}} \frac{\phi}{m_P} \right). \]

Suppose that the Universe is initially dominated by the above modulus. The non-perturbative scalar potential in Eq. (3) is very steep (exponential of an exponential), which means that the field soon becomes dominated by its kinetic density. Once this is so, the particular form of the potential ceases to be of importance. To achieve inflation we assume that, while rolling, the modulus crosses an ESP and becomes temporarily trapped at it.

### 3. At the Enhanced Symmetry Point

In string compactifications there are distinguished points in moduli space at which there is enhancement of the gauge symmetries\textsuperscript{23}. This results in some massive states of the theory becoming massless at these points. Even though from the classical point of view an ESP is not a special point, as the modulus approaches it certain states in the string spectrum

\textsuperscript{a}We considered \( c\sigma > 1 \) to secure the validity of the supergravity approximation and we have assumed that the ESP lies at a minimum in the direction of \( \text{Im}(T) \).
become massless. In turn, these massless modes create an interaction potential that may drive the field back to the symmetry point. In that way a modulus can become trapped at an ESP. The strength of the symmetry point depends on the degree of enhancement of the symmetry.

Such modulus trapping can lead to a period of so-called ‘trapped inflation’, when the trapping is strong enough to make the kinetic density of the modulus fall below the potential density at the ESP. However, it turns out that the number of e-foldings of trapped inflation cannot be very large. Therefore, with respect to cosmology, the main virtue of the ESPs relies on their ability to trap the field and hold it there, at least temporarily.

Because ESPs are fixed points of the symmetries we have

$$V'(\phi_0) = 0,$$

where the prime denotes derivative with respect to $\phi$ and $\phi_0$ is the value of the modulus at the ESP. The above means that the ESP is located either at a local extremum (maximum or minimum) or at a flat inflection point of the scalar potential, where $V'(\phi_0) = V''(\phi_0) = 0$. This means that the presence of an ESP deforms the non-perturbative scalar potential (see Fig. 1). This deformation may be enough so that, after trapped inflation, the field undergoes slow-roll inflation over the flat region of the scalar potential at the vicinity of the ESP. The total duration of inflation may, thus, be enough to solve the flatness and horizon problems of the HBB.

![Fig. 1. Illustration of how the appearance of an ESP at $\phi = \phi_0$ deforms the non-perturbative scalar potential $V_{np}$ to generate, for example, a local maximum at potential density $V_0$. The crossing modulus is temporarily trapped by the emergence of an interaction potential $V_{int}$ due to its enhanced interaction with other fields. After released from trapping, the modulus may drive slow-roll inflation while sliding over the potential hill.](image-url)
3.1. Trapped Inflation

Let us briefly study the trapping of the modulus at the ESP. We assume that around the ESP there is a contribution to the scalar potential due to the enhanced interaction between the modulus $\phi$ and another field $\chi$, which we take to be also a scalar field. The interaction potential is

$$V_{\text{int}}(\phi, \chi) = \frac{1}{2} g^2 \chi^2 \bar{\phi}^2,$$

where $\bar{\phi} \equiv \phi - \phi_0$ with $g$ being a dimensionless coupling constant.

Thus, at the ESP the $\chi$ particles are massless. The time dependence of the effective (mass)$^2$ of the $\chi$ field results in the creation of $\chi$-particles. This takes place when the field is within the production window $|\phi| < \Delta \phi \sim (\dot{\phi}_0/g)^{1/2}$, where $\frac{1}{2} \dot{\phi}_0^2$ is the kinetic density of the modulus when crossing the ESP and the dot denotes derivative with respect to the cosmic time $t$.

The effective scalar potential near the ESP is $V_{\text{eff}}(\phi) \approx V_0 + \frac{1}{2} g^2 \langle \chi^2 \rangle \bar{\phi}^2$ where $V_0 \equiv V(\phi_0)$ with $V(\phi)$ being the 'background' scalar potential. Following Ref.\cite{10} we have $\langle \chi^2 \rangle \simeq n_\chi / g |\phi|$, where $n_\chi$ denotes the number density of $\chi$ particles produced after the crossing of the ESP. This means that $V_{\text{eff}}(\phi) \sim V_0 + gn_\chi |\phi|$ and the field climbs a linear potential since $n_\chi$ is constant outside the production window.

After the first crossing, the field reaches the amplitude $\Phi_1$, determined by its initial kinetic density. To avoid overshooting the ESP we require $\Phi_1 \lesssim m_P$ since for larger values the coupling softens. After reaching $\Phi_1$, the field reverses direction and crosses the production window again, generating more $\chi$ particles and, therefore, increasing $n_\chi$. Thus, it now has to climb a steeper potential reaching an amplitude $\Phi_2 < \Phi_1$. The process continues until the ever decreasing amplitude becomes comparable to the production window (see Fig. 2). At this moment particle production stops.

After the end of particle production, $\langle \chi^2 \rangle$ remains roughly constant during an oscillation and the modulus continues oscillating in the quadratic interaction potential. Studying this oscillation, we found that, due to the Universe expansion, the amplitude and frequency decrease as $\Phi \sim \Delta \phi / a$ and $\langle \chi^2 \rangle \propto a^{-2}$,\cite{9} where the scale factor $a(t)$ is normalised to unity at the end of particle production. Hence, the quadratic potential becomes gradually “diluted” due to the Universe expansion (see Fig. 3). The above mean that the kinetic density of the oscillating modulus scales as $\rho_{\text{osc}} \propto a^{-4}$. When $\rho_{\text{osc}}$ becomes redshifted below $V_0$, trapped inflation begins.

The above process assimilates a multitude of initial conditions (provided overshooting the ESP is avoided) because any kinetic density in excess of $V_0$ is depleted before the onset of trapped inflation. Trapped inflation dilutes
Fig. 2. Illustration of the trapping of a modulus crossing the ESP during particle production. Outside the production window, the modulus oscillates in a linear interaction potential, which steepens progressively due to the production of more $\chi$-particles every time the modulus crosses the ESP.

Fig. 3. Illustration of the trapping of a modulus crossing the ESP after particle production. Inside the production window, the modulus oscillates in a quadratic interaction potential, which becomes gradually diluted due to the Universe expansion.

exponentially the density of the $\chi$–particles, which quickly redshifts $V_{\text{int}}$. Therefore, after a rather limited number of e-foldings of trapped inflation, the modulus is released from the ESP.
3.2. Slow-Roll inflation

Since the ESP is located at a locally flat region of the potential there is a chance that, after $V_{\text{int}}$ becomes negligible, the modulus drives a period of slow-roll inflation while sliding away from the ESP. To study this period we need to quantify the deformation of the scalar potential due to an ESP.

The appearance of an ESP generates either a local extremum or a flat inflection point at $\phi_0$. In all cases, in the vicinity of the ESP, the scalar potential can be approximated by a cubic polynomial. Hence, the characteristics of the potential depend only on $m_\phi^2 \equiv V''(\phi_0)$ and $V^{(3)}_0 \equiv V'''(\phi_0)$. In fact, we can parametrise the deformation of the scalar potential using

$$|V^{(3)}_0| \sim \xi^2 \sigma_0^3 H_*^2/m_P^2,$$

(7)

where $\sigma_0 \equiv \sigma(\phi_0)$ and $H_*$ is the Hubble parameter during inflation: $H_*^2 \approx V_0/3m_P^2$. The $\xi$ parameter accounts for the strength of the symmetry point; the smaller the $\xi$, the stronger the deformation and the wider the inflationary plateau. The requirement that the deformation becomes negligible at distances larger than $m_P$ results in the lower bound $\xi > 1$, which also guarantees that the modulus does not overshoot the ESP.

By studying inflation after the modulus escapes trapping, we have obtained the following results, depending on the ESP morphology.

In each case, one has to achieve enough inflationary e-folds to solve the horizon and flatness problems, while also taking care that the curvature perturbations due to the modulus are not excessive compared to observations.

Consider first the case of a flat inflection point. In this case, we can have enough e-folds of slow roll inflation if $|V^{(3)}_0| < g^2 H_* \ll H_*$. The case of a local minimum is indistinguishable from the above if $m_\phi^2 < g^2 |V^{(3)}_0|$. If the opposite is true then the modulus becomes trapped in the local minimum and must escape through tunnelling. Afterwards the modulus can drive a period of slow-roll inflation with total number of e-foldings given by $N \sim (H_*^2/m_\phi^2)$. Hence, to solve the horizon and flatness problems we need $m_\phi \ll H_*$. Finally, in the case of a local maximum, after the end of trapping, one can have a phase of fast/slow roll inflation provided $|m_\phi| \lesssim H_*$. Thus, we have found that, in all cases, enough slow-roll inflation to solve the horizon and flatness problems of the HBB is attainable provided $|m_\phi|, |V^{(3)}_0| < H_*$. Choosing for illustrative purposes an intermediate value for the Hubble scale: $H_* \sim 1$ TeV we have found that one can achieve enough inflationary e-folds (up to $N_{\text{max}} \sim 10^4$) without producing excessive curvature perturbation if $1 < \xi^2 < 10^4$. Thus, there is ample parameter space for slow-roll inflation to occur after the modulus escapes trapping.
at the ESP. Note also, that, while $H^*$ is determined by the location of the ESP in field space (by the vacuum density $V_0$), the values of $|m|_{V_0}$ and $|V_0(3)|$ are due to the deformation of the scalar potential at the vicinity of the ESP which is not directly related to $V_0$. Hence, the requirement that the latter are smaller than $H^*$ does not necessarily imply fine-tuning.

4. After the end of inflation

After inflation, the field rolls away from the ESP. Soon the influence of the ESP on the scalar potential diminishes and $V(\phi) \approx V_{np}(\phi)$. The steepness of $V_{np}$ results in the kinetic domination of the modulus density. As a result a period of kination occurs, during which the field equation is: $\ddot{\phi} + 3H\dot{\phi} \approx 0$. Hence, the density of the Universe scales as $\rho \propto a^{-6}$. During kination, the scalar field is oblivious of the particular form of the potential.

Kination is terminated when the density of the decay products of a curvaton field dominates the kinetic density of the modulus. Thus, the end of kination corresponds to reheating, with reheating temperature $T_{reh} \sim \sqrt{H_{reh}m_p}$, where $H_{reh}$ is the Hubble parameter at reheating.

After the onset of the HBB, the rolling scalar field is subject to cosmological friction, which asymptotically freezes the field at the value $\phi_F/m_p \simeq \frac{1}{V_0} \ln(V_0/T_{reh}^4)$. Note that this value depends on $T_{reh}$ which, in turn, is determined by curvaton physics. The modulus remains frozen until the present when it plays the role of quintessence. This guarantees that there is no dangerous variation of fundamental constants during the HBB. The evolution of the modulus until today is depicted in Fig. 4.

5. Quintessence

Since $\sigma_F \equiv \sigma(\phi_F) \sim (V_0/T_{reh}^4)^{1/3} > 1$, the modulus rolls to large values before freezing. At such values we can assume that the scalar potential is

$$V(\sigma) \simeq \frac{C_n}{\sigma^n} \Rightarrow V(\phi) \simeq C_n e^{-b \phi/m_p},$$

where $C_n$ is a density scale and $b = \sqrt{2/3} n$. The above is a typical uplift potential introduced by flux compactifications as discussed below.

If the modulus is to account for the required dark energy, it must satisfy the coincidence requirement: $V(\sigma_F) \simeq \Omega_\Lambda \rho_0$, where $\Omega_\Lambda \simeq 0.73$ is the dark energy density parameter and $\rho_0$ is the critical density at present. Hence,

$$T_{reh} \sim V_0^{1/4} (\rho_0/C_n)^{3/8n^2}.$$
Inflation, $\rho_{\text{curv}}$ is subdominant and remains constant until, after the end of inflation (denoted by ‘end’) the curvaton begins oscillating (at time denoted by ‘osc’). During the oscillations, $\rho_{\text{curv}}$ scales as pressureless matter. Sometime afterwards (denoted by ‘dec’) the curvaton decays into the thermal bath of the HBB. This thermal bath dominates the Universe at reheating (denoted by ‘reh’) soon after which the modulus freezes (at time denoted by ‘frz’) assuming constant potential density comparable to the density today.

Thus, the density scale $C_n$ is determined by $T_{\text{reh}}$ which, in turn, is determined by curvaton physics. An upper bound on $C_n$ is obtained by demanding that reheating occurs before big bang nucleosynthesis (BBN):

$$C_n \lesssim \rho_0 \left( \frac{V_0^{1/4}}{T_{\text{BBN}}} \right)^{2n\sqrt{2/3}},$$

where $T_{\text{BBN}} \sim 1$ MeV is the temperature at BBN.

The scalar potential in Eq. (8) may have a multitude of origins. For example, using the volume modulus, we may consider a stack of $D3$–branes...
located at the tip of a Klebanov-Strassler throat. The uplift potential is 
\[ \delta V \sim \exp(-8\pi K/3M g_s) m_P^4/\sigma^2 \equiv C_2/\sigma^2, \] 
where \( M \) and \( K \), in the warp factor, are the units of RR and NS three-form fluxes. To satisfy Eq. (10) we must have \( C_2^{1/4} \lesssim 10^{-20}m_P \). This can be attained by choosing the ratio of fluxes as \( K/M g_s \gtrsim 22 \). Taking \( g_s = 0.1 \), only twice as many units of \( K \) flux as those of \( M \) flux are needed.

It is also possible to consider fluxes of gauge fields on \( D7 \)-branes. In this case, the scalar potential obtains a contribution 
\[ \delta V \sim 2\pi E^2/\sigma^3 \equiv C_3/\sigma^3, \] 
where \( E \) depends on the strength of the gauge fields considered. The constraint in Eq. (10) requires now \( C_3^{1/4} \lesssim 10^{-15}m_P \sim 1 \text{ TeV} \).

The future of the modulus after unfreezing depends on the steepness of the scalar potential, or equivalently the value of \( b \) in Eq. (8).

- For \( b \leq \sqrt{2} \), the modulus dominates the Universe for ever, leading to eternal acceleration. This results in future horizons, which pose a problem for the formulation of the S-matrix in string theory.
- For \( \sqrt{2} < b \leq \sqrt{3} \), the modulus dominates the Universe but results only in a brief accelerated expansion period. Such is the fate of the \( n = 2 \) case.
- For \( \sqrt{3} < b \leq \sqrt{6} \), the modulus does not dominate the Universe, albeit causing a brief period of accelerated expansion. Afterwards the modulus density remains at a constant ratio with the background matter density. This is the fate of the \( n = 3 \) case.
- For \( b > \sqrt{6} \), the modulus does not cause any accelerated expansion and so cannot be used as quintessence. After unfreezing, the modulus rolls fast down the quintessential tail of the scalar potential with its density approaching asymptotically kinetic domination (and subsequently freezing at a value larger than \( \sigma_F \)). This case corresponds to \( n > 3 \).

The brief acceleration period caused by the unfreezing modulus is due to the fact that the modulus oscillates around an attractor solution, which in itself does not result to acceleration. In Ref. it was claimed that brief acceleration occurs if \( \sqrt{2} < b \leq 2\sqrt{6} \), which corresponds to the range \( \sqrt{3} < n \leq 6 \). More recent studies, however, have reduced this range. Brief acceleration in the range \( \sqrt{2} < b \leq \sqrt{3} \) has been confirmed by Ref. This includes the \( n = 2 \) case which corresponds to the most popular uplift potential. The range for \( b \) was expanded further in Ref., where it is shown that brief acceleration can explain the observations at least up to \( b \approx \frac{3}{4} \sqrt{6} \approx 1.837 \). Since the data are interpreted using a number of priors, we believe that the \( n = 3 \) case is still marginally acceptable.
6. Conclusions

Quintessential inflation is possible to achieve in string theory with flux compactifications using as inflaton a string modulus which rolls down its runaway potential. Inflation is due to the presence of an enhanced symmetry point (ESP), which traps the modulus and creates a locally flat region over which the modulus can slow-roll. There is ample parameter space for successful inflation provided: \( m_\phi, |V^{(3)}_0| \ll H_* \). Trapping assimilates a multitude of initial conditions provided overshooting the ESP is avoided.

After inflation, the modulus becomes (again) kinetically dominated causing a period of kination. Reheating is due to the Universe domination by the decay products of a curvaton field, which also accounts for the correct amplitude of the curvature perturbations in the Universe. During the Hot Big Bang, the modulus freezes and remains constant until the present.

At the frozen value, the potential is dominated by an uplift term of exponential form. This residual potential density begins to dominate today, when the modulus unfreezes, leading to a brief period of late inflation.

Curvaton physics fixes the reheating temperature and determines the value of the frozen modulus \( \sigma_F \) and the density scale \( C_n \) in the uplift potential. Coincidence and BBN constrains on \( C_n \) allow realistic values much less fine-tunned than the cosmological constant \( \Lambda \) in the \( \Lambda \)CDM model.

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