Introduction: Inflation,[1][4], a period of exponential expansion driven by a scalar field called the inflaton, is the current paradigm for the origin of our universe, and in particular for the seeds of structure within it that eventually grew into clusters, galaxies, stars and planets. Inflation speaks to several theoretical puzzles: why is the universe spatially flat? Why does the universe look homogeneous and isotropic and the largest scales, seemingly implying equilibration of causally disconnected regions? But most importantly, inflationary models make clear predictions [5–10] for the spectrum of density fluctuations in the early universe. Quantum fluctuations get stretched to macroscopic scales by the inflationary expansion, and gravity does the rest. Overdense regions attract more matter which eventually collapses under its own attractions and leads to all the visible structure in the universe. The almost scale invariant spectrum implied by inflation is in stunning agreement with observations of the cosmic microwave background [11].

Most inflationary models require the evolution of the scalar field φ to be “slow roll”. That is the scalar field is slowly rolling down a relatively shallow scalar potential. Technically this is incorporated in the so called slow roll parameters which characterize the slope and curvature of the potential $V(\phi)$:

$$\epsilon \equiv \frac{M_P^2}{16\pi} \left(\frac{V'}{V}\right)^2, \quad \eta \equiv \frac{M_P^2}{8\pi} \left(\frac{V''}{V}\right).$$

(1)

Slow roll inflation corresponds to choosing $\epsilon, \eta \ll 1$. In all simple inflationary models slow roll is absolutely necessary for two reasons. For one, the number of e-folds, that is the logarithm of the final over initial scale factor of the universe, is proportional to $\epsilon^{-1/2}$. To get a sufficient number of e-folds to solve the horizon and flatness problems to begin with one needs either a small epsilon or large changes in the scalar field value compared to the Planck scale. More quantitatively, $\epsilon$ and $\eta$ lead to deviations from a scale invariant power spectrum of the fluctuations and so the experimentally observed almost scale invariant spectrum directly implies upper bounds on $\epsilon$ and $\eta$. This is quantified via the so called spectral tilt $n_s - 1$. $n_s = 1$ corresponds to a completely scale invariant power spectrum, whereas $n_s$ that deviates from that value tell us that the power spectrum has a small variation with $k$. In a slow roll inflation model one can find that [12]

$$n_s - 1 = -6\epsilon + 2\eta. \quad (2)$$

The degeneracy between $\epsilon$ and $\eta$ can be broken by considering higher order variations in $k$. Using the 2018 data from the Planck mission on the observed power spectrum in the cosmic microwave background (CMB) radiation one can derive an observational upper bound on $\epsilon$ of

$$\epsilon < 0.0063 \quad (3)$$

at the 95% confidence level [11]. However, potentials support the smallness of $\epsilon$ are generally unnatural with the same reason as the long-standing issue of Higgs mass in the Standard Model of particle physics. Furthermore, inflation also has other deep theoretical shortcomings, see e.g. [13], that question some of its phenomenological successes. Those problems are almost all about the unknown underlying origin of inflation namely looking at it from an effective field theory point of view without addressing its ultra-violet(UV) completion.

This urges a better understanding of reality from an underlying UV complete point of view. String theory is the only known consistent framework of a UV complete unifying theory and so a lot of string theory motivated conjectures, called swampland conjectures, distinguishing those effective field theories able to be embedded in string theory to those cannot, are proposed [14][16]. Among these conjectures, a so called de-Sitter swampland conjecture [15] and its refined version [17] strongly
denied the existence of slow roll inflation \[17, 18\]. However, the potential addressed by those authors should be clarified. As from several explicit examples provided by \[15\] from string theory compactifications, we here claim that the potential addressed by them is the geometric potential which is part of Einstein-Hilbert action from a higher dimensional decompactified point of view. Hence these potentials did not fully address the interaction between the scalar field and hidden sectors.

In this work, we will try to study the whole quantum interaction effect of string theory compactification onto inflation, due to the so called distance conjecture \[16\], and see that in a complete description of inflation, with the de-Sitter conjecture and the quantum interaction effect, slow roll inflation actually has a potential mechanism to be engineered from string theory. Therefore, the existing problems \[13\] of inflation might have a unified answer in this complete description. We’ll discuss issues about fine-tuning at the end.

**Setup and notation:** We will consider a scalar field constructed from a consistent quantum gravitational theory, for example string theory. To address the complete dynamics of this scalar field including potential interactions with other sectors, we will use the distance conjecture \[14\].

**Distance Conjecture:** As the modulus moves a distance beyond Plank scale in moduli space, there would be a tower of light states emerges with masses exponentially suppressed.

In this statement, modulus refers to a scalar field \(\phi\) parametrizing the compactification or roughly speaking the size of an extra dimension and, in our context, we will take it to be the inflaton. The emergent modes are usually the so called Klauzau-Klein (KK) modes and winding modes from string theory compactification \[19, 20\] (other examples can be found at \[14\]). Moreover, the KK modes will be heavy for tiny extra dimensions and will be light for large extra dimensions and the winding modes will behave in a opposite way. Without loss of generality, we will consider a single copy of the emergent mode described as a scalar field \(\Phi\) with mass square

\[
m^2 = m^2 e^{-c \frac{\phi}{M_P}}
\]

where \(M_P\) is the Plank scale, \(c\) is an order one constant and \(m_0^2\) is defined in a way that \(\phi\) starts at 0.

From now on we will call the modulus field \(\phi\) as inflaton. We will consider the following Lagrangian density

\[
\mathcal{L} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) + \frac{1}{2} \partial_\mu \Phi \partial^\mu \Phi - \frac{1}{2} m^2(\phi) \Phi^2
\]

where \(\Phi\) is the emergent mode from the distance conjecture and \(V(\phi)\) is the potential of inflaton from string theory compactification. This potential satisfies the following conjecture \[15\].

**De-Sitter Swampland Conjecture:** Scalar field potentials arising from a consistent quantum gravitational theory should satisfy

\[
|\nabla V| \geq \mathring{c} V
\]

where \(\mathring{c}\) is of order one in Plank unit.

As it is usually done \[21\], we will take the exponential form of the inflaton potential

\[
V(\phi) = B e^{-\frac{\phi}{m_0}}, \quad b \sim O(1).
\]

The comment on inflation and current experiment on the CMB power spectrum \[3\] from this conjecture is that it is in tension with inflation but not with \[3\]. Because as we discussed in the introduction that to resolve the horizon and flatness problems using inflation we need slow roll \(\epsilon, \eta \ll 1\). And to satisfy the current experimental upper bound we only have to require that \(b < 0.564\) which is not necessarily in conflict with \(b \sim O(1)\).

**Integrating out the emergent modes:** Since we want to identify a mechanism for slow roll inflation, we’ll assume \(\phi\) is slow rolling and check that if there is a mechanism supporting this assumption given that the dS swampland conjecture is satisfied. We want to consider the complete dynamics of \(\phi\) including its interaction with the emergent mode and hence we will integrate out the emergent mode. Also, we assume that the characteristic mass scale \(m_0\) of the emergent modes is heavier than that of \(\phi\). As a prelude, we’ll exactly integrate it out and focus on the resulting dynamics for the inflation. The result is that we have a quantum effective theory of inflaton organized systematically by the small expansion parameter \(\frac{\phi}{m_0}\) and \(O(x) = e^{-\frac{\phi}{m_0}} - 1\) and they are small because of the slow roll of \(\phi\). Our resulting effective potential is

\[
V_{\text{eff}} = \mathcal{B} e^{-\frac{\phi}{m_0}} + \frac{m_0^4}{32\pi^2} \left[ \log^2 \left( \frac{\mu^2}{m_0^2} \right) (1 - e^{-2\frac{\phi}{m_0}}) + \frac{1}{2} \log \left( \frac{\mu^2 e^2}{m_0^2} \right) (1 - e^{-2\frac{\phi}{m_0}})^2 + \frac{c_0^2 e^{-4\frac{\phi}{m_0}}}{M_{pl}} \right]
\]

where \(\mu\) is the cutoff scale beyond which a complementary light state will take over. In the language of string theory we can take \(\mu\) to be the self-dual scale under T-duality \[19, 20, 22\]. Details of how to derive this exact effective potential can be found in the appendix.

**Slow roll parameters:** Slow roll parameters are defined in \[14\]. In this section and next we will see that slow roll can be supported by our effective potential \[8\]. We will use the fact that among the slow roll regime, \(\phi\) starting from the origin, \(e^{-\frac{\phi}{m_0}} \sim 1\). The slow roll parameters are given approximately by

\[
\epsilon \approx \frac{b^2}{16\pi} \left[ 1 - \frac{m_0^2}{16\pi^2 B b} \frac{\log^2 \left( \frac{\mu^2}{m_0^2} \right)}{m_0^2} \right],
\]

where
\[ \eta \approx \frac{b^2}{8\pi} \left[ 1 + \frac{m_0^4}{4\pi^2 B b^2 M_p} \phi (1 - 2 \log(\frac{\mu^2}{m_0^2})) \right]. \tag{10} \]

**Discussion and future remarks:** From the expressions of slow roll parameters \([11, 12]\), we see that the cutoff dependence is logarithmic and hence our theory is not fine-tuned from the field theory point of view. However, in some sense we do need a "fine-tuning" by adjusting several parameters for example \(B, b, c, m_0\) and the self-dual scale \(\mu\) but they can be engineered in string theory compactification. As a result, from this UV-complete description we saw that fine-tuning is not a technical problem for our mechanism of inflation but just a way to look at the huge landscape of string theory vacua. A possible question is that for string theory compactification we always have KK modes and winding modes the masses of which behave in opposite ways as the size of the extra dimension changes so in addition to \([13]\) there should be other modes with exponentially growing mass. And these heavy modes would have opposite contributions to slow roll parameters from the light modes. However, the point is that distance conjecture is true only for trans-Planckian moduli distances namely when the extra dimension parametrized by \(\phi\) is either very close to zero or very large. In the intermediate regime these two effects are competing and hence cancel each other in the loop. And one of them will dominate over the other for trans-Planckian moduli distances which is the case we are considering. Another possible question about how large \(\frac{\mu}{m_0}\) is because it determines whether the logarithmic terms are positive or negative and how positive or negative they are. We could say for sure that \(m_0\) is lower than the self-dual scale \(\mu\) because as we said before that if it goes over \(\mu\) then the complementary sector will be light and take over in the low energy effective theory. As a result, it is for sure that the logarithmic terms are positive and hence \(\epsilon\) is reduced. Now the question is if it is possible that the logarithmic terms are bigger than one and therefore \(\eta\) will also be reduced. For this question we did not see any obstacle why this is not possible. And this the reason why we think that our mechanism is a potential mechanism but not a literal mechanism for inflation. If it is the case that the logarithmic term must be smaller than one then we might have to consider the refined de-Sitter swampland conjecture \([10]\). A potential short coming of our mechanism is that we did not consider the field theory in a curved spacetime. Moreover, our effective action can be used to understand the relationship between infinite distance and an infinite tower of emergent modes from the point of effective field theory as suggested in \([14]\) in a closed form to all orders in perturbation theory. We will leave these questions for future studies.

The last thing we would like to address is that the geometry of string theory might change our way of thinking about low energy physics from the traditional quantum field theory point of view (more examples of this view point can be reached at e.g. \([22, 24]\)).

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**Details on the effective action**

In this section we will show details of our effective theory which is one loop exact. Our effective theory is controlled by two small parameters \(\frac{\epsilon}{\mu^2}\) and \(\mathcal{O}(k) = (e^{-\frac{\epsilon}{\mu^2}} - 1)(-k)\). The first one is small because we are in slow roll regime which is relevant for inflation and the second one is small in the sense that as an operator its expectation value is small because we defined our theory near \(\phi = 0\) and slow roll.

The effective action of \(\phi\) reads

\[
S_{\text{eff}} = S_{\text{free}} + \frac{i}{2} \text{Tr} \log[\partial^2 + m_0^2] e^{-\frac{2\phi}{M_p} \partial}\]

\[
= S_{\text{free}} + \frac{i}{2} \text{Tr} \log[(\partial^2 + m_0^2)(1 + \frac{m_0^2(e^{-\frac{2\phi}{M_p}} - 1)}{\partial^2 + m_0^2})]
\]

\[
= S_{\text{free}} + \frac{i}{2} \text{Tr} \log[1 + \frac{m_0^2(e^{-\frac{2\phi}{M_p}} - 1)}{\partial^2 + m_0^2}]
\]

\[
= S_{\text{free}} + \frac{i}{2} \text{Tr} \log[1 + \frac{1}{1 + \frac{m_0^2}{m_0^2}} \mathcal{O}(k)]
\]

\[
= S_{\text{free}} + \frac{i}{2} \text{Tr} \log[1 + \frac{1}{1 + \frac{m_0^2}{m_0^2}} \mathcal{O}(k)]
\]

\[
(11)
\]

where the \(\log(\partial^2 + m_0^2)\) in the second step only gives an overall normalization factor of the partition function and hence is dropped.

From here we see that we can systematically reorganize the effective action using the two small parameters mentioned at the beginning of this section. Furthermore, this is the exact effective action of the inflaton (the only Feynman diagrams are of type \([11]\), defined at energy scales below \(m_0\) and the regime that \(\phi\) is slowly rolling. The effective potential, i.e. the non-derivative part of the effective action is found to be \([8]\). \(\mu\) is the cutoff scale of the emergent modes inside the loop and it is the scale beyond which the complementary emergent mode will take over. For example, if it is a KK mode then the complementary mode is a winding mode and vice versa.

The first term in \(\mathcal{O}(k)\) expansion can be worked out as
A. A. Starobinsky, A. H. Guth, where potential we can drop the derivative terms. Furthermore, Derivative starts to appear at the second term

\[
\Gamma^{(2)} = - \frac{i}{4} \text{Tr} \left[ \frac{m_0^2}{\partial^2 + m_0^2} \mathcal{O}(x) \frac{m_0^2}{\partial^2 + m_0^2} \mathcal{O}(x) \right]
\]

(12)

where \( \eta \) is the Feynman parameter and for the effective potential we can drop the derivative terms. Furthermore, we can work out the general term for \( n > 2 \) as

\[
\Gamma^{(n)} = \frac{i(-1)^{n-1}}{2n} \text{Tr} \left\{ \left[ \frac{m_0^2}{\partial^2 + m_0^2} \mathcal{O} \right]^n \right\}
\]

\[
= - \frac{i m_0^4}{2n} \int \frac{d^4k}{(2\pi)^4} \frac{d^4p}{(2\pi)^4} \frac{1}{\partial^2 + m_0^2} \frac{1}{\partial^2 + m_0^2} \times \tilde{\mathcal{O}}(k) \tilde{\mathcal{O}}(-k)
\]

\[
= \frac{m_0^4}{64\pi^2} \int \frac{d^4k}{(2\pi)^4} \int_0^1 d\eta \log \left( \frac{\tilde{\mu}^2}{m_0^2 - \eta(1-\eta)k^2} \right) \times \tilde{\mathcal{O}}(k) \tilde{\mathcal{O}}(-k)
\]

\[
= \frac{m_0^4}{64\pi^2} \int \frac{d^4x}{64\pi^2} \int_0^1 d\eta \mathcal{O}(x) \log \left( \frac{\tilde{\mu}^2}{m_0^2 + \eta(1-\eta)\partial^2} \right) \mathcal{O}(x)
\]

(13)

Derivative starts to appear at the second term

\[
\Gamma^{(1)} = \frac{i}{2} \text{Tr} \left[ \frac{m_0^2}{\partial^2 + m_0^2} \mathcal{O}(x) \right]
\]

\[
= \frac{i}{2} \int d^4x \mathcal{O}(x) \int \frac{d^4k}{(2\pi)^4} \frac{m_0^2}{-k^2 + m_0^2}
\]

\[
= - \frac{1}{32\pi^2} \int d^4x \mathcal{O}(x) (m_0^4)^2 \log \left( \frac{\tilde{\mu}^2}{m_0^2} \right).
\]

If we drop the terms vanishing in the limit \( m_0^2 \ll \tilde{\mu}^2 \) then we get

\[
\Gamma^n = \frac{(-1)^n m_0^4}{32\pi^2 n(n-1)(n-2)} \int dx_1 \cdots dx_n \frac{d^4k_1}{(2\pi)^4} \cdots \frac{d^4k_n}{(2\pi)^4} \tilde{\mathcal{O}}(k_1) \cdots \tilde{\mathcal{O}}(k_n) \delta \left( \sum_{i=1}^n x_i - 1 \right) \left( \frac{\Delta}{m_0^2} + 1 \right)^{-n/2}
\]

(14)

where \( \Delta \) is from Feynman parametrization given as

\[
\Delta = \left[ x_2 k_2 + x_3 (k_2 + k_3) + \cdots + x_n (k_2 + \cdots + k_n) \right]^2
\]

\[
- x_2 k_2^2 - \cdots - x_n (k_2 + \cdots + k_n)^2.
\]

(15)

To extract the zeroth order term in \( \frac{\tilde{\mu}^2}{m_0^2} \) expansion, we replace \( \left( \frac{\Delta}{m_0^2} + 1 \right)^{-n/2} \) by 1. The zeroth order term reads

\[
\Gamma_0^n = \frac{(-1)^n m_0^4}{32\pi^2 n(n-1)(n-2)} \int dx \frac{d^4k_1}{(2\pi)^4} \cdots \frac{d^4k_n}{(2\pi)^4} \tilde{\mathcal{O}}(k_1) \cdots \tilde{\mathcal{O}}(k_n)
\]

\[
\times (2\pi)^4 \delta^{(4)}(k_1 + \cdots + k_n) \times (2\pi)^4 \delta^{(4)}(k_1 + \cdots + k_n)
\]

\[
\times \cdots \left[ (p + k_2 + \cdots + k_n)^2 - m_0^2 \right]
\]

(17)

Summing over all these contributions we get the desired result [5].

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