Eur. Phys. J. Spec. Top.
https://doi.org/10.1140/epjs/s11734-021-00195-w

THE EUROPEAN PHYSICAL JOURNAL SPECIAL TOPICS

Review

Models of accelerating universe in supergravity and string theory

Koushik Dutta[1,a] and Anshuman Maharana[2,b]

[1] Indian Institute of Science Education and Research Kolkata, Mohanpur, WB 741246, India
[2] Harish-Chandra Research Institute, HBNI, Allahabad 211019, India

Received 2 March 2021 / Accepted 4 June 2021

Abstract We review models of the accelerating universe from the perspective of high-energy physics. Focusing on supergravity and the String Theory, we discuss the general framework for the construction of these models. We then go on to discuss explicit constructions.

1 Introduction

The discovery of a transition from a decelerating to an accelerating Universe in the recent epoch of cosmic evolution is one of the greatest scientific achievements of modern times. The observational data suggest that around 70% of the total energy density of the current Universe belongs to an unknown component of the Universe, broadly called dark energy.[1] The observations suggest that the current energy density of dark energy is $\rho_{\rm DE} \sim 10^{-47}$ GeV^4. More specifically, the equation of state parameter of the dark energy fluid is reported to be $w_0 = -1.028 \pm 0.032$ [1].

The simplest solution for the dark energy is the existence of a positive cosmological constant Λ with the constant equation of state parameter $w = -1$. The positive contributions to the cosmological constant come from the vacuum fluctuations of any quantum field, and its estimated value is typically much larger than the observed value. We pretend this contribution to the cosmological constant is zero, or at the least, it does not gravitate [2]. On the other hand, the dark energy can also be described by the value of a scalar field potential at its minimum $V(\phi_{\rm min})$, also called vacuum energy, whose magnitude at present is approximate $V(\phi_{\rm min}) \sim \rho_{\rm DE} \sim \Lambda^4$. The other possibility that is often considered is quintessence where the vacuum energy can be positive, zero, or negative, while at present the quintessence field is slow-rolling at positive value of its potential with $V(\phi_{\rm today}) \sim \rho_{\rm DE}$. Crucially, in this case, the w is time-dependent. If the future observations find $w \neq -1$, it will be crucial to understand the fundamental origin of the field, and it will confirm the existence of new dynamical matter d.o.f in the Universe [3].

In the context of vacuum energy, the important question is to find scalar field potential with a de Sitter minimum. This is an issue that is discussed in the context of supergravity, or more aptly in the context of String Theory. De Sitter constructions remain a subject of intense debate, particularly in light of the swampland conjectures on dS initiated with the papers[2] [4,5]. In this review article, we shall not discuss de Sitter vacua via examining the swampland conjecture, progress in proving/disproving them and their implications. The literature on swampland conjectures is vast, we refer the reader to the review [9] for a comprehensive discussion of the subject. Our approach will be to give a summary of the de Sitter constructions in string theory and also a description of further work required to put them in a firm footing.

In an effective field theory approach, the value of $V(\phi_{\rm min})$ is obtained (in principle) by integrating out all modes from the UV to the cosmological constant scale Λ. In string models, we integrate out all string modes and Kaluza–Klein modes down to the compactification scale $M_{\rm KK}$ to find out the 4D effective action. It gives the vacuum energy $\langle V \rangle$ at scale $M_{\rm KK}$ which is different from $V(\phi_{min})$ as it does not take into account corrections coming from integrating out light degrees of freedom associated with any energy scale $\tilde{M}$ between $M_{\rm KK}$ and Λ. Therefore, generically we will have $V(\phi_{min}) = \langle V \rangle + \mathcal{O}(\tilde{M}^4)$. For models where supersymmetry (SUSY) is used for explanation of the gauge hierarchy, one expects contributions $\tilde{M} \sim$ TeV, otherwise, even larger energy scales could also be present. Thus,

[1] The observed acceleration may also happen due to the modifications of gravity at cosmological distances.

[a] e-mail: koushik.physics@gmail.com (corresponding author)
[b] e-mail: anshumanmaharana@hri.res.in

[2] For earlier criticisms of dS in string theory, see, for example, [6–8].

it is important to be able to obtain four-dimensional string vacua with a value of the vacuum energy which can be tuned to cancel low-energy $\mathcal{O}(\tilde{M}^4)$ corrections leading to $V(\phi_{min}) \sim \Lambda^4$. This is usually argued to be possible due to the presence of the flux landscape where one can choose background fluxes so that $\langle V \rangle$ cancels off the low energy correction. Note that similar arguments hold for the constructions of the quintessence potential which has appropriate magnitude, but with the additional constraint that the mass for the scalar field must be tuned as well as be protected from quantum corrections to make it cosmologically relevant.

This article is structured in the following manner: in the next section, we outline the basic requirements for the quintessence field with a brief outlines of typical potential forms. In Sect. 3, we discuss quintessence models in the framework of supergravity which includes the effects of SUSY breaking as well as quintessence being a pNGB. We then proceed to discuss the issue in the context of String Theory.

2 Requirement for a quintessence field

In the simplest form of scalar field dark energy model, it is assumed that a canonically normalized scalar field ϕ, called quintessence, is minimally coupled with gravity, and it dominates the energy density of the Universe at the current epoch [10–13]. In accordance with the observation, if $\dot{\phi}^2 < V(\phi)$, the Universe accelerates, but the stringent requirement of $w \simeq -1$ enforces the field to roll very slowly, *i.e* $\dot{\phi}^2 \ll V(\phi)$. Moreover, a substantial amount of dark matter at the present epoch plays a crucial role in the dynamics of quintessence. The fact that the quintessence field is away from its minimum and dynamical at the present epoch requires the field to be extremely light, namely its mass $m_\phi \lesssim H_0 \sim 10^{-33}$ eV, where H_0 is the present Hubble constant. Finally, related to the quintessence potential, there must be a mass scale $\sim 10^{-3}$ eV which essentially determines the current energy density contributions from the field. The fundamental origin of this scale along with the value of m_ϕ is a crucial issue to be understood. These two fine-tuning of parameters exist even at the classical level of a given potential.

From the phenomenological point of view, there are many potentials that can satisfy the above requirements. Broadly, those potentials can be discussed in two categories [14]. The first one being the thawing models where the field was frozen due to the Hubble damping at the earlier epochs, and only started to roll in recent time. A well-known example of this kind of model is the periodic potential. For the case of the freezing model, the field gradually slows down as the Universe evolves. Among all the potentials, the inverse power-law potentials of the form

$$V(\phi) = M^{4+\alpha}/\phi^\alpha, \tag{1}$$

with $\alpha > 0$ requires special attention due to its tracking behaviour of the background energy density [11–13]. In this case, both the coincidence problem and the initial problem have natural resolutions.

Unless there is a special symmetry, the quintessence field typically has $\mathcal{O}(1)$ couplings with the Standard Model particles which will lead to a long-range fifth force of gravitational strength, as well as it introduces time variations of the fundamental constants [15]. The observations put severe constraints on these couplings [16,17]. Even with $\mathcal{O}(1)$ couplings, these constraints can be evaded by several mechanisms [18,19]. On the other hand, if the quintessence field is a pseudo-Nambu–Goldstone–Boson (pNGB), i.e a pseudo-scalar these couplings can be suppressed due to their derivative nature [15].

It is straight forward to write down effective Lagrangian with potential $V_0(\phi)$ for the quintessence field that satisfies all the above conditions. In the context of a more fundamental theory, the crucial point is the natural origin of the dark energy scale and the radiative stability of the light mass of the field. In addition to the quintessence field ϕ, if a theory involves a field with a mass $\tilde{M}$, once the heavy field is integrated out, the quintessence potential typically receives corrections proportional to $\tilde{M}^4$. These corrections vanish in the presence of unbroken SUSY, but SUSY is broken at a much higher scale than Λ. In effect, the corrections are $\gtrsim \mathrm{TeV}^4$, and it spoils the flatness of the original potential V_0 that we started with. This issue will be much clearer in coming sections where we discuss quintessence potential in the context of supergravity and SUSY breaking.

On the other hand, if the field originates from a spontaneously broken global symmetry, it is a pNGB with a periodic potential. In this case, radiative stability of the potential is guaranteed by the associated shift symmetry [20]. On the other hand, the observations require the symmetry breaking scale $f \gtrsim M_{\rm pl}$ [21,22], where $M_{\rm pl}$ is the reduced Planck mass. This is difficult to accommodate in the context of a fundamental theory like String Theory [23,24]. As we will see, a pNGB like axion with its suppressed couplings with the SM particles and with radiative stability of its mass, it is a promising candidate for the quintessence field.

3 Models in supergravity

The nature of the quintessence field as described in the previous section does not allow it to be identified with a Standard Model (SM) d.o.f. Therefore, we must look beyond the SM physics to identify the field. Among several possibilities, SUSY is being considered the most promising candidate. For typical quintessence potential, the field value at the present epoch is of order the Planck mass, and it enforces us to consider the effects of gravity in a supersymmetric theory, i.e. in supergravity (SUGRA). The supergravity theories can be thought of as the effective four-dimensional action of the String Theory where the low-energy effective action can be derived where heavy fields have been integrated out. We now discuss certain representative models of quintessence in supergravity, and in the end, we discuss the generic difficulties with these models.

3.1 Power-law models

Before discussing models in supergravity, let us review the first model of supersymmetric quintessence. Using nonperturbative effects in globally supersymmetric gauge theories, the inverse power law quintessence potential of the form of Eq. (1) has been constructed in [25,26] using the following superpotential:

$$W = \Lambda^{3+\gamma}\phi^{-\gamma}, \tag{2}$$

where $2(1+\gamma) = \alpha$ and $\Lambda = M$. But, this set-up is excluded by the current observational data as w always remain larger than -0.7. On theoretical side, the typical field value is larger than the UV cut-off scale Λ. Moreover, the analysis does not incorporate the supergravity corrections when the field tracks at the present epoch with $\phi \sim M_{\rm pl}$.

In $\mathcal{N} = 1$ supergravity the F-term scalar potential (we ignore the D-term contributions to the potential for simplicity now) for a chiral field φ is given in terms of Kahler potential K, and the superpotential W

$$V = e^{K/M_{\rm pl}^2}\left(K^{ij^*}D_iW(D_jW)^* - 3|W|^2/M_{\rm pl}^2\right)\ , \tag{3}$$

where $D_iW \equiv \partial W/\partial\varphi^i + (W/M_{\rm pl})\partial K/\partial\varphi^i$, and $K_{ij^*} \equiv \partial^2 K/\partial\varphi^i\partial\varphi^{j*}$ with K^{ij*} is inverse of K_{ij*}. Although the last term in the bracket allows us to realise SUSY breaking ($D_iW \neq 0$) with zero vacuum energy ($V = 0$), on the other hand it puts obstacle in realising positive vacuum energy. Using the same superpotential as above in Eq. (2) and $K = \phi\phi^*$, the potential looks like (setting $M_{\rm pl} = 1$) [27,28]

$$V = e^{(\phi^2/2)}\frac{\Lambda^{4+\beta}}{\phi^\beta}\left[(\beta-2)^2/4 - (\beta+1)(\phi^2/2) + \phi^4/4\right], \tag{4}$$

where $\beta = 2\gamma + 2$. Note that the the potential becomes negative for $\phi \sim M_{\rm pl}$. Even though this problem can be solved by choosing particular values of Λ, but that value does not satisfy other observational constraints. This is true for all values of γ.

The above mentioned general problem in supergravity can be solved either by choosing a specific superpotential with the property of its expectation value $\langle W\rangle = 0$ [27,28] (but its derivative being non-zero), or more general Kahler potential [29] than the canonical one we have chosen till now. For the first case, theory consists of charged matter field Y in addition to the quintessence field. The field rolls along a D-flat direction with D-term potential contributions being zero, and F-term of a matter field contributes to the potential energy such that the potential becomes positive [27,28]

$$V = e^{K/M_{pl}^2}K^{YY^*}|W_Y|^2\ , \tag{5}$$

and the potential becomes of the form

$$V = e^{\frac{\phi^2}{2}}\frac{\Lambda^{4+\gamma}}{\phi^\gamma}. \tag{6}$$

First, note that the potential is positive, and an important difference appear due to the exponential factor. The equation of state parameter w is pushed toward further lower values $w \simeq -0.86$ in comparison to the usual inverse power law potential of Eq. (1), and this result is nearly independent of the values of γ. Still, note that the model is excluded by the current observational data [1].

The negativity of supergravity potential of Eq. (3) can also be taken care by the appropriate choice of the Kahler potential. With the choice of superpotential $W = \Lambda^{3+\gamma}\tilde{\phi}^{-\gamma}$ of the form of Eq. (2), and the Kahler potential of the form $K = -M_{pl}^2\ ln[(\tilde{\phi}+\tilde{\phi}^*)/M_{pl}]$, the scalar potential in terms of canonically normalised quintessence field ϕ looks like [29]

$$V = M^4 e^{-\sqrt{2}\beta\phi/M_{pl}}\ , \tag{7}$$

where $M^4 = M_{pl}^{-\beta-1}\Lambda^{\beta+5}(\beta^2-3)/2$, and $\beta = 2\gamma+1$. For $\beta > \sqrt{3}$ the potential is positive, and a scaling solution exists where the ratio of quintessence energy density and the matter energy density remain constant. Obviously, this solution can not make transition to a quintessence dominated era, and thus the potential needs to be appropriately modified. The reference [29] further proposed a Kahler potential of the form $K = M_{pl}^2[ln\ (\tilde{\phi}+\tilde{\phi}^*)/M_{pl}]^2$. For $|\phi| \ll M_{pl}$, the potential behaves as $V \sim (-\phi)^{-2/3}$, and for $|\phi| \gg M_{pl}$, $V \sim (-\phi)^{-2/3}e^{(-\phi/M_{pl})^{4/3}}$ with having a minimum with positive potential energy in the intermediate field range $\phi \sim M_{pl}$. In this case, depending on the initial values of the field two different kinds of attractor solutions exist. In both the cases, when the field approaches its minimum it starts to dominate the energy density of the Universe, and the field starts behaving as a cosmological constant with $w \simeq -1$. Broadly, in this category of models, the dynamics of the field are non-trivial due to non-canonical terms, and therefore, a potential that might look unsuitable for dark energy becomes consistent in terms of the canonically normalized field.

3.2 SUSY breaking effects and other issues

The above models in pure supergravity look fine, even though not fully satisfactory in terms of recent stringent observational constraints. In addition to that it is important to remember that in reality the SUSY must be broken. It turns out that the effect of SUSY breaking generically induces unacceptably large corrections to the otherwise suitable quintessence potential in supergravity [30].

The F-term scalar potential can be written in the following form

$$V_F = |F|^2 - 3m_{3/2}^2 M_{pl}^2, \tag{8}$$

where the non-zero vev of $F = e^{K/M_{\rm pl}} K^{ij*} D_i W (D_j W)^*$ corresponds to the SUSY breaking, and the gravitino mass $m_{3/2} = e^{K/2M_{\rm pl}^2} |W|/M_{pl}^2$. Considering the potential V_F must be of the order of dark energy scale, we need $F \sim m_{3/2} M_{pl}$ which essentially corresponds to a corrections $\delta V \sim m_{3/2}^2 M_{pl}^2$ to the supergravity potential due to SUSY breaking. Now, $\sqrt{F} \sim 10^{10}$GeV, and 10^4GeV for gravity- and gauge-mediated SUSY breaking respectively, and it makes $\delta V \gg V_F$, and it spoils the flatness of the original potential.

This can be explicitly seen with with the following separable Kahler potential $K = \phi\phi^* + XX^*$ and the superpotential $W = W(\phi) + \tilde{W}(X)$, where X is the chiral field responsible for SUSY breaking.[3] In this case, the F-term potential will contain a term which is proportional to $|W|^2|\phi|^2 \sim m_{3/2}^2 |\phi|^2$ such that

$$V_{tot} = V_0 + m_{3/2}^2 |\phi|^2. \tag{9}$$

Here V_0 is the potential without SUSY breaking effects being taken care of. Thus, the SUSY breaking effects induce a mass to the quintessence field of the order of $m_{3/2}$ which is much heavier than the required mass for acceleration to drive. In this case, the field quickly rolls to its minimum and behaves as a nondynamical cosmological constant depending on the value of the potential at its minimum.

The solution to this generic problem requires an unconventional mechanism of SUSY breaking. For example, due to the effects of higher-dimensional physics, it is possible that even though mass differences between the supersymmetric particles arise, the SUSY may still remain unbroken, and thus does not contribute any corrections to the potential. In this case, the quintessence sector and the SUSY breaking sector is sequestered [38]. In some extended models of supergravity, the mass of the scalar field is not related to the SUSY breaking scale. In this case $m_\phi^2 = nH_0^2$ where n is an integer. For example, in the $\mathcal{N} = 2$ case $n = 6$, and $\mathcal{N} = 8$ case $n = -6$ with the following potential [39,40]

$$V = 3H_0^2 M_{pl}^2 \left(1 \pm (\phi/M_{pl})^2\right) . \tag{10}$$

In the context of supergravity, another generic problem is the couplings of the quintessence field to the supersymmetric version of the SM fields (MSSM). These couplings induce the time dependence of fermion masses, violations of the weak equivalence principle, variations of the gauge coupling constants, and extra fifth force carried by the field [41,42].

Many of these problems can be resolved by demanding shift symmetry in the Kahler potential of the quintessence field [43]. Due to the underlying symmetry, higher-order shift symmetry breaking corrections to the Kahler potential does not make large contributions to the potential. For the case of [43] the dynamics of the field becomes indistinguishable from the cosmological constant, whereas for [38] the potential become axionic—see the next section.

3.3 pNGB models

A crucial approach for dark energy model building is the idea that the quintessence field is a pseudo-Nambu–Goldstone–Boson (pNGB) [20]. In this case, the underlying theory posses a $U(1)$ global symmetry which gets spontaneously broken at the scale f, and an angular field ϕ appears as a massless Goldstone Boson. Due to an explicit symmetry breaking at the scale μ, the scalar field acquires a mass, and the potential looks in the following form:

$$V(\phi) = \mu^4 \left(1 + Cos(\phi/f)\right). \tag{11}$$

In the limit $\mu \to 0$, the potential vanishes. The radiative corrections to the potential are proportional to μ^4 due to the underlying symmetry, and therefore, its mass is technically natural and protected from possible quantum corrections. At the same time, the approximate symmetry also suppresses the couplings between ϕ and the matter fields [15]. The most famous example of the pNGB is the QCD axion where the underlying $U(1)$ is the Peccei–Quinn symmetry.

When the above potential drives cosmic acceleration, we need $\mu \sim 10^{-3}$eV, and $f \gtrsim M_{pl}$ [21,22]. In this case, the field starts rolling along with the potential very recently when $m_\phi \sim H_0$, and thus behaving as a thawing dark energy model. In the context of high-energy physics, it is crucial to understand the origin of the scale μ. In reference [31], the quintessence field is a model-independent axion field whose potential scale originates from the hidden sector gaugino condensation. It mixes with another hidden sector axion where one linear combination behaves as dark energy and the other one behaves as the dark matter related to the QCD-axion. The idea of generating the quintessence potential from electroweak instanton effects when axion couples to electroweak gauge fields have been explored in [32]. In ref [33], the scale of the axion potential is related to the electroweak scale v by see-saw like mechanism where the dark energy scale $\mu \sim v^2/M_{pl}$, and $m_\phi \propto v^4/M_{pl}^3$. For several other ideas where

[3] Note that in more realistic cases, this separation may not be possible, making the problem much worse.

a pNGB is the quintessence field, see the references [34–36].

4 De Sitter in string theory for dark energy

Moduli stabilisation is best understood in the context of type II B models, therefore we will focus mostly on type II B. An attractive feature of the IIB models is the no-scale structure, this keeps the Kähler moduli massless at tree level. These directions are then stabilised by the inclusion of perturbative (in both α' and g_s) and non-perturbative corrections to the effective action, leading to stabilisation at weak coupling and large volume. In this way, one can hope to address the the Dine–Seiberg problem [44]. For this reason, several dS mechanisms have been proposed and developed in IIB. We first discuss II B models in detail and then go on to briefly discuss models in other settings. For more technical summaries of the state of the art in dS constructions see for example [45, 46].

4.1 Overview of IIB flux compactifications

Type IIB compactifications on orientifolds of a Calabi–Yau (CY) manifolds have many features that make them attractive for phenomenology. We briefly review the structure of their low-energy effective field theory (EFT). The relevant closed string fields are the the complex structure moduli U_a, $a = 1, \cdots, h^{1,2}$ axio-dilaton S, and the Kähler moduli T_i, $i = 1, \cdots, h^{1,1}$ where $h^{1,2}$ and $h^{1,1}$ are the Hodge numbers of the compact CY manifold (after orientifolding). The tree-level Kähler potential is given by :

$$K = -2\ln\mathcal{V} - \ln\left(S+\bar{S}\right) - \ln\left(-\mathrm{i}\int_X \Omega\wedge\bar{\Omega}\right), \tag{12}$$

where $\mathcal{V} = \ell_s^{-6}\int_X \sqrt{g_{(6)}}\, d^6y$ is the volume of the internal manifold in units of the string length ℓ_s. The internal volume $\mathcal{V}$ is a homogeneous function of degree 3/2 of the real parts of the Kähler moduli τ_i that determine the sizes of four cycles. Ω is the holomorphic (3, 0)-form of the manifold. In the presence of fluxes, the superpotential is given by [47]:

$$W_{\text{flux}} = \int_X G_3\wedge\Omega\,, \tag{13}$$

where the 3-form flux $G_3 = F_3 - \mathrm{i}SH_3$. These 3-form fluxes are quantised. The superpotential (13) thus fixes the dilaton and all complex structure moduli and reduces the number of vacua from a continuum to a discrete but large set of points [48, 49]. The equations of motion require G_3 to be imaginary self-dual, i.e. $*_6G_3 = \mathrm{i}G_3$; in terms of Hodge decomposition $G_3 \in (2,1)\oplus(0,3)$. Note that in general the solutions break supersymmetry (supersymmetry is preserved only if the $(0,3)$ component is turned off).

The Kähler moduli are not stabilised by the fluxes. This is due to the fact that there exists a Peccei–Quinn symmetry $T_i \to T_i + \mathrm{i}c_i$ with constant c_i's that together with the holomorphy of the superpotential, does not allow any T_i dependence of W to all orders in perturbation theory. However these moduli are the gauge couplings for matter fields on D7-branes, and so effects like gaugino condensation on D7-branes or Euclidean D3-instantons (see e.g [50]) generate a non-perturbative superpotential for them. The full superpotential for closed string moduli is

$$W = W_{\text{flux}}(S,U) + W_{\text{np}}(S,U,T)\,. \tag{14}$$

The starting point for the search for vacua in the 4D EFT is the F-term supergravity scalar potential for arbitrary superpotential $W(\Phi_M)$ and Kähler potential $K(\Phi_M, \bar{\Phi}_{\bar{M}})$ (3), which we record again for ease in the discussion. In units of $M_{\rm p}$,

$$V_F = e^K\left(K^{M\overline{N}}\,D_MW\overline{D}_{\overline{M}}\overline{W} - 3|W|^2\right), \tag{15}$$

where $D_MW = \partial_MW + (\partial_MK)W$. The tree-level Kähler potential for the Kähler moduli satisfies the no-scale condition $K^{T_i\bar{T}_{\bar{j}}}K_{T_i}K_{\bar{T}_{\bar{j}}} = 3$. Two major scenarios have emerged to fix the Kähler moduli: the KKLT proposal [51] and the Large Volume Scenario (LVS) [52–54]. Both are in the $W_0 \neq 0$ case. KKLT makes use of the fact that W_0 can be tuned to small values and then compete with the small non-perturbative effects in $W_{\rm np}$ to produce an AdS minimum for the T-fields. In this case, the minimum is at $D_{T_i}W = 0$ and is supersymmetric. For LVS, instead of tuning W_0, the leading order no-scale breaking effect, which is a $\mathcal{V}$-dependent α' correction, competes with non-perturbative corrections on a small (blow-up) four cycle. At the AdS minimum obtained, the volume $\mathcal{V} \sim e^{1/g_s} \gg 1$ is exponentially large in string units and supersymmetry is broken by the F-terms associated with the Kähler moduli.

Since perturbative and non-perturbative effects play a central role in fixing the Kähler moduli, we begin by sketching the general structure of these corrections to K and W. First, since string theory has no free parameter, each coupling corresponds to the value of a different modulus: the string coupling $g_s = 1/\mathrm{Re}(S)$ is set by the dilaton, the Kähler moduli control α' effects and the coupling of gauge theories on D7-branes wrapping internal four cycles. Thus moduli stabilising corresponds to fixing the values of the expansion parameters. Contrary to field theories, string compactifications have feature many expansion parameters. This makes it hard to obtain exact results but also provides more flexibility regarding weak coupling expansions. At weak coupling, the leading order correction to the tree-level Kähler potential for the T-moduli in (12) arises from perturbative effects (either in α' or g_s) and we denote this as

$$K = -2\ln\mathcal{V} + K_{\rm p}\,. \tag{16}$$

We will write superpotential (14) in the schematic form:

$$W = W_0 + W_{\rm np}\,. \tag{17}$$

The F-term scalar potential can be expanded as

$$V = V_0 + \delta V\,, \tag{18}$$

where the tree-level potential V_0 is positive definite due to the no-scale symmetry. Since V_0 is independent of the Kahler moduli, the minimum of the potential in the Kähler moduli space is determined by the corrections δV. Thus, determining the leading contributions to δV is central to properly stabilising the moduli. From the expressions in (16) and (17), the structure of δV takes the form [53]:

$$\delta V \propto e^K \left(W_0^2\, K_{\rm p} + W_0\, W_{\rm np}\right)\,. \tag{19}$$

If there were only one expansion parameter, and if $W_0 \gg W_{\rm np}$ and $K_{\rm p} \gg W_{\rm np}$ the the first term would be the leading order term. It would lift the potential, but would necessarily give rise to a runaway behaviour, unless terms of different order in the perturbative expansion compete to give a minimum. This can however arise only in a regime where the perturbative expansion would break down since the corresponding expansion parameter would not be small. This is the essence of the Dine–Seiberg problem [44].

Type IIB flux compactifications, however, provide two ways to overcome this issue. First, in the KKLT scenario the big discrete degeneracy of flux vacua is used to tune W_0 to be exponentially small such that $W_0 \sim W_{\rm np}$. This then requires that the $W_{\rm np}^2$ terms to be also included in (19), they stabilise the T-fields as they compete with the $W_0\, W_{\rm np}$ terms [51]. Note that in this regime quantum corrections to the Kähler potential can be consistently neglected since the first term in (19) is subdominant since $W_0^2\, K_{\rm p} \ll W_0, W_{\rm np} \sim W_0^2$ for $K_{\rm p} \ll 1$.

The second possibility is exploited in the LVS models. Here, the fact that there is more than one expansion parameter plays the key rôle. In this case, the two terms in (19) can compete with each other to provide a minimum as long as each arises from a different expansion. Thus, at the minimum one has $W_0^2\, K_{\rm p} \sim W_0\, W_{\rm np}$ which, for $K_{\rm p} \sim 1/\mathcal{V}$ and $W_{\rm np} \sim e^{-\tau_s}$, yields an overall volume of order $\mathcal{V} \sim W_0\, e^{\tau_s}$. Here τ_s is a blow-up four cycle that gets stabilised at a value of order of $1/g_s$. It is therefore large at weak string coupling, as a result the CY volume is exponentially large [52–54].

In summary, KKLT exploits tuning of the fluxes to obtain $W_0 \sim W_{\rm np} \ll 1$. Recently, a systematic procedure to obtain vacua with low W_0 has been put forward in [55]. Whereas LVS works for natural values of the flux superpotential of order $W_0 \sim \mathcal{O}(1-100)$ (see [56] for a concrete example) but depends more on perturbative corrections to K.

4.1.1 Advantages

Let emphasise several advantages of type IIB constructions:

1. *Suppressed scalar potential scale*: The starting point of the dS models is the classical low-energy effective action of type IIB string theory compactified on a CY orientifold. This is a controlled approximation if the compactification volume is large and the energies involved (E) are such that

$$E \ll M_{\rm KK} = \frac{M_s}{\mathcal{V}^{1/6}} \ll M_s \equiv \frac{1}{\ell_s} \equiv \frac{1}{2\pi\sqrt{\alpha'}} = g_s^{1/4}\frac{M_{\rm P}}{\sqrt{4\pi\mathcal{V}}}\,, \tag{20}$$

   where M_s is the string scale. As mentioned above, the scale of the potential at tree-level is of order $V_0 \sim M_s^4$ but the vacuum energy vanishes due to the no-scale cancellation. This allows one to keep the value of the scalar potential around the minimum well below the string and the Kaluza–Klein scale, and the effective field theory approach justified.
2. *Controlled flux backreaction*: Fluxes can be turned on to generate a potential for the moduli in a controlled fashion since their backreaction on the internal geometry just makes the compactification manifold conformally Calabi–Yau. Thus, the understanding of the underlying moduli space is better than in other settings. Progress has been made recently in computing the form of the Kähler potential including the effects of warping [58–65]. We note that the warping induces corrections to the definition of the correct moduli coordinates which are however negligible at large volume.
3. *Suppressed SUSY breaking scale*: Supersymmetry is broken at tree-level by the F-terms of the Kähler moduli which are proportional to the $(0,3)$ component of G_3 and scale as $F^T = e^{K/2}\, K^{T\bar{T}} K_{\bar{T}} W_0 \sim \frac{W_0}{\mathcal{V}^{1/3}}$. Thus the scale of supersymmetry breaking is well below the KK scale: the gravitino mass $m_{3/2} = e^{K/2} W_0 \sim \frac{W_0}{\mathcal{V}}$ is hierarchically smaller than the Kaluza–Klein scale $M_{\rm KK} \sim M_s/\mathcal{V}^{1/6} \sim 1/\mathcal{V}^{2/3}$ for either $W_0 \ll 1$ (as in KKLT constructions) or $\mathcal{V} \gg 1$ (as in LVS models).[4] For a recent analysis of the distribution of the scale of supersymmetry breaking in IIB models, see [66].
4. *Progress in computing quantum effects*: Much progress has been made in computing non-perturbative contributions to the superpotential (Euclidean D3-brane instantons in particular [50]) and perturbative corrections to the Kähler potential. After the first computation of $N = 2$ $\mathcal{O}(\alpha'^3)$

[4] Note that in F-theory models the string coupling can be arbitrarily large, this tree-level analysis is not valid.

corrections to the Kähler potential K [67], additional $N = 2$ $\mathcal{O}(g_s^2\alpha'^2)$ and $\mathcal{O}(g_s^2\alpha'^4)$ contributions to K have been obtained in [68] and further advanced in [69]. Ref. [70] showed the existence of an extended no-scale structure since $\mathcal{O}(g_s^2\alpha'^2)$ contributions to the scalar potential vanish. More recently, there has been substantial progress in understanding $N = 1$ perturbative effects. Reference [71] showed that $N = 1$ $\mathcal{O}(\alpha'^2)$ corrections to the effective action lead to moduli redefinitions, while Ref. [72] found that $N = 1$ $\mathcal{O}(\alpha'^3)$ effects are captured by a shift of the CY Euler number term.[5] Moreover, Ref. [74] reconsidered $N = 2$ $\mathcal{O}(\alpha'^3)$ contributions to the Kahler potential incorporating the backreaction of these corrections on the internal geometry and found that they lead to moduli redefinitions. Progress has also been made in the computation of higher derivative $N = 2$ $\mathcal{O}(\alpha'^3)$ terms [75,76] which can have implications for moduli stabilisation and cosmology [77,78]. Finally Ref. [79–81] have recently obtained $N = 1$ string loop corrections to the Einstein-Hilbert term showing that they generate g_s^2 corrections to a term involving the CY Euler number.[6]
It is worth emphasising that none of the perturbative α' and g_s corrections listed above create instabilities for LVS constructions. On the other hand, corrections subleading in an inverse volume expansion turn out to be useful to lift leading order flat directions. We also note it is at times not necessary to obtain the full functional dependence of these corrections on all moduli, but it is sufficient to determine their dependence on the Kähler moduli which are yet to be stabilised. The functional dependence of string loop corrections to K on the Kähler moduli can often be determined from generalisations of toroidal computations and low-energy arguments [70,83]. Another powerful requirement is the positivity of the Kähler metric (see section 5.2 of [84]). It has been shown in [85], that even though the flux superpotential W_0 depends explicitly on the dilaton it is still not renormalised at any order in perturbation theory. This is non-trivial since the standard arguments for the non-renormalisability of W used the fact that W did not depend on the string coupling [86–88].
5. *Controlled higher derivative corrections*: As shown in [89], the superspace derivative expansion is justified if $W_0 \ll \mathcal{V}^{1/3}$ which corresponds to requiring a gravitino mass which much smaller than the Kaluza–Klein scale. As discussed earlier, this can be achieved by either tuning $W_0 \ll 1$ as in KKLT models or by $\mathcal{V} \gg 1$ as in LVS constructions.
6. *dS mechanisms*: Several mechanisms have been proposed to obtain dS vacua in type IIB models. Some of the better examined dS mechanisms are: (i) anti-branes [51], (ii) T-branes [90], (iii) α' effects [91], (iv) non-perturbative effects at singularities [92], (v) non-zero S and U F-terms [93].
7. *Explicit global models*: A full model, should not only lead to a dS vacuum but also include SM-like chiral matter and an inflationary sector. A lot of progress in this direction has been made in the type IIB [57, 97–104]. Supersymmetry breaking in such settings has been studied in detail recently in [105,106].

4.2 Other dS mechanisms

Some other string motivated proposals are:

- *Non-critical strings* [107]: Non-critical strings have a positive cosmological term (for $D > 10$) that can be used to obtain dS by compactification while fixing the moduli. The major challenge here is obtaining a better understanding of the EFT, so as to establish that the approximations are under control.
- *Negative curvature spaces* [108]: Non-supersymmetric compactifications on manifolds with negative curvature lead to a positive term in the effective potential that can yield dS solutions. Being non-supersymmetric, the EFT is under less control but these compactifications can have interesting implications for the landscape.
- *Kähler uplift* [91,109,110]: The α' corrections to the Kähler potential in the KKLT scenario can be made to compete with the fluxes and the non-perturbative effects to produce solutions with positive vacuum energy. The dS minima so obtained are in regions at the edge of validity of the EFT. An explicit construction has been carried out in [56] with all geometric moduli stabilised.
- *Dilaton-dependent non-perturbative effects* [92]: In type IIB models hidden and observable sectors can be localised in D3 or D7-branes. One possibility is to consider dilaton-dependent non-perturbative effects coming from $E(-1)$-instantons or strong dynamics on a hidden sector of D3-branes at singularities. The non-perturbative term yields a positive definite contribution to the scalar potential similar to that which arises from anti-branes.
- *Complex structure F-terms* [93]: The complex structure moduli have minima at the supersymmetric points $D_U W = 0$. But there may be other minima for these fields for which $D_U W \neq 0$. These can in principle give rise to dS minima without the need of further ingredients.. A concrete example with $\mathcal{V} \simeq 10^4$ was constructed in [93].
- *Non-perturbative dS vacua* [94]: dS minima can emerge from stabilising all the geometric moduli in just one-step via the inclusion of just background fluxes and non-perturbative effects. The main problems of this approach are: the poor knowledge of the S and U-moduli dependence of the prefactor of non-perturbative effects and the computational difficulty to find a numerical solution for the minimi-

[5] See also [73] for $N = 1$ $\mathcal{O}(\alpha'^2)$ corrections to K in heterotic strings which should get mapped to type IIB $\mathcal{O}(g_s^2\alpha'^2)$ effects that have the extended no-scale cancellation.
[6] See also [82].

sation equations in the presence of a large number of geometric moduli.

- *Heterotic dS vacua* [95,96]: The heterotic strings on smooth Calabi–Yau threefolds yield dS models where the gauge bundle moduli, together with the dilaton and the U-moduli are fixed supersymmetrically at leading order. The Kähler moduli can be fixed as in LVS by the interplay of worldsheet instantons, α' effects and threshold corrections to the gauge kinetic function.
- *G_2 compactifications* [111,112]: G_2 holonomy compactifications of M-theory can also lead to dS vacua. Here, interesting scenarios have been studied addressing phenomenological issues. Superpotentials with two exponentials but without fluxes have been proposed to stabilise the moduli. The freedom to tune the cosmological constant arises by scanning through different ranks of the condensing gauge groups.
- *IIA models* Type IIA models allow for fixing of all the moduli at tree-level by background fluxes. However, so far no stable dS vacua have been found [113–119]. These constructions have the advantage of stabilising all the moduli at the classical level. But the ten-dimensional equations can be solved exactly only under the approximation of smearing which would lead to a Calabi–Yau internal manifold. But, in the localised case, the four dimensional effective field theory is not under control since the backreaction of the fluxes on the internal manifold cannot be neglected leading to a half-flat non-Calabi–Yau metric [120,121].
- *Non-Geometric Constructions* Non-geometric constructions seem to yield dS vacua without tachyons [122–126]. However the exact form of the moduli space is unknown, this is required to check the validity of the approximations. Moreover, the tree-level stabilisation procedure leads to a four-dimensional potential of order the string scale with $\mathcal{O}(1)$ values of the volume of the internal manifold, thus it is not clear if α' effects can consistently be ignored.

5 Quintessence in string theory

The key challenges for quintessence in string theory are the bounds from fifth forces and finding potentials which are flat after quantum corrections are incorporated (for a recent discussion of these issues, see, e.g [127]). Axions are ideal candidates to overcome these difficulties are their shift symmetry protects their potential at the perturbative level. Furthermore, they couple to matter through derivative couplings; as a result the fifth force constraints are not severe. Thus they have been widely used for constructing models of quintessence in String Theory [128–135]. In fact, a generic prediction of string compactifications is the string axiverse [136]. A large number of axions arising upon dimensional reduction of the antisymmetric form fields on the internal cycles of the compactification space. The number of axions is given by to the number of cycles in the compactification manifold which can easily be of O(100) or larger. Shift symmetries prevent them from acquiring masses at the perturbative level. The mass of an axionic particle is given by

$$m_a \sim M_{\rm pl} e^{-\tau_a}, \tag{21}$$

where τ_a is the volume (in string units) of the cycle on which the non-perturbative effect responsible for the mass of the axion is supported. One can have compactification scenarios in which a good number of the τ_a are large, leading to a many ultra-light axions. The potential for such axions takes the form

$$V = \Lambda^4 - \sum_{i=1}^{N_{\rm ULA}} \Lambda_i^4 \cos\left(\frac{a_i}{f_i}\right), \tag{22}$$

where $N_{\rm ULA}$ is the number of ultra-light axions. The axions can have a range of masses, the late time cosmological dynamics can be examined by integrating out the heavier ones (with each integrating out, cosmological constant term Λ has to suitably corrected). The potential of Eq. (22) allows for realising various scenarios of quintessence: (a) natural inflation (b) hilltop (c) quasi-natural and (d) oscillating scalar. We refer the reader to [46] and references there in for the details of these scenarios.[7] The detailed phenomenological implications of these scenarios have been explored in [139].

6 Conclusions

The current acceleration of the Universe due to either vacuum energy at the bottom of a scalar potential with positive potential energy or a rolling field (quintessence) which is away from its minimum are well-established mechanisms. The difficulties arise when the idea is incorporated within our knowledge of high-energy physics. The main difficulty arises due to the small scale of the associated energy density in comparison to the known scales of high-energy physics. String compactifications seem to have the necessary ingredients for construction of de Sitter vacua. There has been substantial progress in this direction, but still many open questions remain. These are being actively investigated, particularly in light of the swampland conjectures. In the context of quintessence, the associated difficulties also arise from its ultralight mass and negligible coupling to matter, in particular to baryons. When the idea of quintessence is thought in supergravity, the main difficulty is controlling the SUSY breaking effects to its mass, and avoiding generic couplings to the matter particles. From that point of view, a pNGB is a

[7] Once challenge is to the field range to be suitably large [23,24]. Various possible solutions have been proposed, see for e.g [9,46,137,138] and the references there in.

good starting point that naturally has a small coupling to other fields, and its mass is also technically small. Therefore, as we have seen in this article, a pNGB (or many) is the favourite choice of quintessence model building. Although, getting a working model in string theory under full computational control which evades the bounds and has the suitable field range remains a challenge. Thus, there is a lot to do from the theoretical perspective. At the same time, existing observations have constrained the departure from a pure de Sitter solution today severely.

Acknowledgements KD is partially supported in part by the Grant MTR/2019/000395 and Indo-Russian project Grant DST/INT/RUS/RSF/P-21, both funded by the DST, Govt of India. AM is supported in part by the SERB, DST, Government of India, by the Grant MTR/2019/000267.

References

1. N. Aghanim *et al.* [Planck Collaboration], “Planck 2018 results. VI. Cosmological parameters,” arXiv:1807.06209 [astro-ph.CO]
2. J. Martin, Everything you always wanted to know about the cosmological constant problem (but were afraid to ask). CR Phys. **13**, 566–665 (2012). [arXiv:1205.3365 [astro-ph.CO]]
3. E. J. Copeland, M. Sami, S. Tsujikawa, Dynamics of dark energy. Int. J. Mod. Phys. D **15**, 1753–1936 (2006) [arXiv:hep-th/0603057 [hep-th]]
4. H. Ooguri, C. Vafa, On the geometry of the string landscape and the swampland. Nucl. Phys. B **766**, 21 (2007). [hep-th/0605264]
5. G. Obied, H. Ooguri, L. Spodyneiko, C. Vafa, De Sitter Space and the Swampland. arXiv:1806.08362 [hep-th]
6. T. Banks, The top 10^{500} reasons not to believe in the landscape. arXiv:1208.5715 [hep-th]
7. S. Sethi, Supersymmetry breaking by fluxes. arXiv:1709.03554 [hep-th]
8. I. Bena, M. Grana, N. Halmagyi, On the existence of meta-stable vacua in Klebanov-Strassler. JHEP **1009**, 087 (2010). [arXiv:0912.3519 [hep-th]]
9. E. Palti, Fortsch. Phys. **67** (2019) no.6, 1900037 https://doi.org/10.1002/prop.201900037[arXiv:1903.06239 [hep-th]]
10. C. Wetterich, Cosmology and the Fate of dilatation symmetry. Nucl. Phys. B **302**, 668–696 (1988). [arXiv:1711.03844 [hep-th]]
11. B. Ratra, P.J.E. Peebles, Cosmological consequences of a rolling homogeneous scalar field. Phys. Rev. D **37**, 3406 (1988)
12. R. R. Caldwell, R. Dave, P. J. Steinhardt, Cosmological imprint of an energy component with general equation of state. Phys. Rev. Lett. **80**, 1582–1585 (1998) [arXiv:astro-ph/9708069 [astro-ph]]
13. I. Zlatev, L. M. Wang, P. J. Steinhardt, Quintessence, cosmic coincidence, and the cosmological constant. Phys. Rev. Lett. **82**, 896–899 (1999) [arXiv:astro-ph/9807002 [astro-ph]]
14. R. R. Caldwell, E. V. Linder, The Limits of quintessence. Phys. Rev. Lett. **95**, 141301 (2005) [arXiv:astro-ph/0505494 [astro-ph]]
15. S. M. Carroll, Quintessence and the rest of the world. Phys. Rev. Lett. **81**, 3067–3070 (1998) [arXiv:astro-ph/9806099 [astro-ph]]
16. E. G. Adelberger, B. R. Heckel, A. E. Nelson, Tests of the gravitational inverse square law. Ann. Rev. Nucl. Part. Sci. **53**, 77–121 (2003) [arXiv:hep-ph/0307284 [hep-ph]]
17. C. J. A. P. Martins, The status of varying constants: a review of the physics, searches and implications. [arXiv:1709.02923 [astro-ph.CO]]
18. J. Khoury, A. Weltman, Chameleon fields: awaiting surprises for tests of gravity in space. Phys. Rev. Lett. **93**, 171104 (2004) [arXiv:astro-ph/0309300 [astro-ph]]
19. K. Hinterbichler, J. Khoury, Symmetron fields: screening long-range forces through local symmetry restoration. Phys. Rev. Lett. **104**, 231301 (2010). [arXiv:1001.4525 [hep-th]]
20. J. A. Frieman, C. T. Hill, A. Stebbins, I. Waga, Cosmology with ultralight pseudo Nambu-Goldstone bosons. Phys. Rev. Lett. **75**, 2077–2080 (1995) [arXiv:astro-ph/9505060 [astro-ph]]
21. K. Dutta, L. Sorbo, Confronting pNGB quintessence with data. Phys. Rev. D **75**, 063514 (2007) [arXiv:astro-ph/0612457 [astro-ph]]
22. D. Adak, K. Dutta, Viable dark energy models using pseudo-Nambu-Goldstone bosons. Phys. Rev. D **90**(4), 043502 (2014) [arXiv:1404.1570 [astro-ph.CO]]
23. T. Banks, M. Dine, P. J. Fox, E. Gorbatov, On the possibility of large axion decay constants. JCAP **06**, 001 (2003) [arXiv:hep-th/0303252 [hep-th]]
24. N. Arkani-Hamed, L. Motl, A. Nicolis, C. Vafa, JHEP **06** (2007), 060 https://doi.org/10.1088/1126-6708/2007/06/060[arXiv:hep-th/0601001 [hep-th]]
25. P. Binetruy, Models of dynamical supersymmetry breaking and quintessence. Phys. Rev. D **60**, 063502 (1999) [arXiv:hep-ph/9810553 [hep-ph]]
26. A. Masiero, M. Pietroni, F. Rosati, SUSY QCD and quintessence. Phys. Rev. D **61**, 023504 (2000) [arXiv:hep-ph/9905346 [hep-ph]]
27. P. Brax, J. Martin, Quintessence and supergravity. Phys. Lett. B **468**, 40–45 (1999) [arXiv:astro-ph/9905040 [astro-ph]]
28. P. Brax, J. Martin, The Robustness of quintessence. Phys. Rev. D **61**, 103502 (2000) [arXiv:astro-ph/9912046 [astro-ph]]
29. E.J. Copeland, N.J. Nunes, F. Rosati, Phys. Rev. D **62**, 123503 (2000). [arXiv:hep-ph/0005222 [hep-ph]]
30. C. F. Kolda, D. H. Lyth, Quintessential difficulties. Phys. Lett. B **458**, 197–201 (1999) [arXiv:hep-ph/9811375 [hep-ph]]
31. J. E. Kim, H. P. Nilles, Phys. Lett. B **553**, 1–6 (2003) [arXiv:hep-ph/0210402 [hep-ph]]
32. Y. Nomura, T. Watari, T. Yanagida, Phys. Lett. B **484**, 103–111 (2000) [arXiv:hep-ph/0004182 [hep-ph]]
33. L.J. Hall, Y. Nomura, S.J. Oliver, Phys. Rev. Lett. **95**, 141302 (2005). [arXiv:astro-ph/0503706 [astro-ph]]
34. Z. Chacko, L.J. Hall, Y. Nomura, JCAP **10**, 011 (2004). [arXiv:astro-ph/0405596 [astro-ph]]

35. R. Barbieri, L.J. Hall, S.J. Oliver, A. Strumia, Phys. Lett. B **625**, 189–195 (2005). [arXiv:hep-ph/0505124 [hep-ph]]
36. P.Q. Hung, Nucl. Phys. B **747**, 55–87 (2006). [arXiv:hep-ph/0512282 [hep-ph]]
37. K. Choi, Phys. Rev. D **62**, 043509 (2000). [arXiv:hep-ph/9902292 [hep-ph]]
38. C. I. Chiang, H. Murayama, [arXiv:1808.02279 [hep-th]]
39. R. Kallosh, A.D. Linde, S. Prokushkin, M. Shmakova, Phys. Rev. D **66**, 123503 (2002). [arXiv:hep-th/0208156 [hep-th]]
40. P. Fre, M. Trigiante, A. Van Proeyen, Class. Quant. Grav. **19**, 4167–4194 (2002). [arXiv:hep-th/0205119 [hep-th]]
41. P. Brax, J. Martin, Phys. Rev. D **75**, 083507 (2007). [arXiv:hep-th/0605228 [hep-th]]
42. P. Brax, J. Martin, JCAP **11**, 008 (2006). [arXiv:astro-ph/0606306 [astro-ph]]
43. P. Brax, C. van de Bruck, J. Martin, A.C. Davis, JCAP **09**, 032 (2009). [arXiv:0904.3471 [hep-th]]
44. M. Dine, N. Seiberg, Is the superstring weakly coupled? Phys. Lett. **162B**, 299 (1985)
45. S. Kachru, S.P. Trivedi, Fortsch. Phys. **67**(1–2), 1800086 (2019). https://doi.org/10.1002/prop.201800086 [arXiv:1808.08971 [hep-th]]
46. M. Cicoli, S. De. Alwis, A. Maharana, F. Muia, F. Quevedo, Fortsch. Phys. **67**(1–2), 1800079 (2019). https://doi.org/10.1002/prop.201800079 [arXiv:1808.08967 [hep-th]]
47. S. Gukov, C. Vafa, E. Witten, CFT's from Calabi-Yau fourfolds. Nucl. Phys. B **584** (2000) 69 Erratum: [Nucl. Phys. B **608** (2001) 477] [hep-th/9906070]
48. K. Dasgupta, G. Rajesh, S. Sethi, M theory, orientifolds and G-flux. JHEP **9908**, 023 (1999). [hep-th/9908088]
49. S.B. Giddings, S. Kachru, J. Polchinski, Hierarchies from fluxes in string compactifications. Phys. Rev. D **66**, 106006 (2002). [hep-th/0105097]
50. R. Blumenhagen, M. Cvetic, S. Kachru, T. Weigand, D-Brane instantons in type II orientifolds. Ann. Rev. Nucl. Part. Sci. **59**, 269 (2009). [arXiv:0902.3251 [hep-th]]
51. S. Kachru, R. Kallosh, A.D. Linde, S.P. Trivedi, De Sitter vacua in string theory. Phys. Rev. D **68**, 046005 (2003). [hep-th/0301240]
52. V. Balasubramanian, P. Berglund, J.P. Conlon, F. Quevedo, Systematics of moduli stabilisation in Calabi-Yau flux compactifications. JHEP **0503**, 007 (2005). [hep-th/0502058]
53. J.P. Conlon, F. Quevedo, K. Suruliz, Large-volume flux compactifications: moduli spectrum and D3/D7 soft supersymmetry breaking. JHEP **0508**, 007 (2005). [hep-th/0505076]
54. M. Cicoli, J.P. Conlon, F. Quevedo, General analysis of LARGE volume scenarios with string loop moduli stabilisation. JHEP **0810**, 105 (2008). [arXiv:0805.1029 [hep-th]]
55. M. Demirtas, M. Kim, L. Mcallister, J. Moritz, Phys. Rev. Lett. **124**(21), 211603 (2020). https://doi.org/10.1103/PhysRevLett.124.211603 [arXiv:1912.10047 [hep-th]]
56. J. Louis, M. Rummel, R. Valandro, A. Westphal, Building an explicit de Sitter. JHEP **1210**, 163 (2012). [arXiv:1208.3208 [hep-th]]
57. M. Cicoli, D. Klevers, S. Krippendorf, C. Mayrhofer, F. Quevedo, R. Valandro, Explicit de Sitter flux vacua for global string models with chiral matter. JHEP **1405**, 001 (2014). [arXiv:1312.0014 [hep-th]]
58. S.P. de Alwis, On potentials from fluxes. Phys. Rev. D **68**, 126001 (2003). [hep-th/0307084]
59. S.B. Giddings, A. Maharana, Dynamics of warped compactifications and the shape of the warped landscape. Phys. Rev. D **73**, 126003 (2006). [hep-th/0507158]
60. A.R. Frey, A. Maharana, JHEP **08**, 021 (2006). https://doi.org/10.1088/1126-6708/2006/08/021 [arXiv:hep-th/0603233 [hep-th]]
61. C.P. Burgess, P.G. Camara, S.P. de Alwis, S.B. Giddings, A. Maharana, F. Quevedo, K. Suruliz, JHEP **04**, 053 (2008). https://doi.org/10.1088/1126-6708/2008/04/053 [arXiv:hep-th/0610255 [hep-th]]
62. G. Shiu, G. Torroba, B. Underwood, M.R. Douglas, Dynamics of warped flux compactifications. JHEP **0806**, 024 (2008). [arXiv:0803.3068 [hep-th]]
63. A.R. Frey, G. Torroba, B. Underwood, M.R. Douglas, The universal Kahler modulus in warped compactifications. JHEP **0901**, 036 (2009). [arXiv:0810.5768 [hep-th]]
64. L. Martucci, On moduli and effective theory of $N = 1$ warped flux compactifications. JHEP **0905**, 027 (2009). [arXiv:0902.4031 [hep-th]]
65. L. Martucci, Warping the Kahler potential of F-theory/IIB flux compactifications. JHEP **1503**, 067 (2015). [arXiv:1411.2623 [hep-th]]
66. I. Broeckel, M. Cicoli, A. Maharana, K. Singh, K. Sinha, JHEP **10**, 015 (2020). https://doi.org/10.1007/JHEP10(2020)015 [arXiv:2007.04327 [hep-th]]
67. K. Becker, M. Becker, M. Haack, J. Louis, Supersymmetry breaking and alpha-prime corrections to flux induced potentials. JHEP **0206**, 060 (2002). [hep-th/0204254]
68. M. Berg, M. Haack, B. Kors, String loop corrections to Kahler potentials in orientifolds. JHEP **0511**, 030 (2005). [hep-th/0508043]
69. M. Berg, M. Haack, E. Pajer, Jumping through loops: on soft terms from large volume compactifications. JHEP **0709**, 031 (2007). [arXiv:0704.0737 [hep-th]]
70. M. Cicoli, J.P. Conlon, F. Quevedo, Systematics of string loop corrections in type IIB Calabi-Yau flux compactifications. JHEP **0801**, 052 (2008). [arXiv:0708.1873 [hep-th]]
71. T.W. Grimm, R. Savelli, M. Weissenbacher, On α' corrections in $N = 1$ F-theory compactifications. Phys. Lett. B **725**, 431 (2013). [arXiv:1303.3317 [hep-th]]
72. R. Minasian, T.G. Pugh, R. Savelli, F-theory at order α'^3. JHEP **1510**, 050 (2015). [arXiv:1506.06756 [hep-th]]
73. L. Anguelova, C. Quigley, S. Sethi, The leading quantum corrections to stringy Kahler potentials. JHEP **1010**, 065 (2010). [arXiv:1007.4793 [hep-th]]
74. F. Bonetti, M. Weissenbacher, The Euler characteristic correction to the Kahler potential—revisited. JHEP **1701**, 003 (2017). [arXiv:1608.01300 [hep-th]]

75. D. Ciupke, J. Louis, A. Westphal, Higher-derivative supergravity and Moduli stabilization. JHEP **1510**, 094 (2015). [arXiv:1505.03092 [hep-th]]
76. T.W. Grimm, K. Mayer, M. Weissenbacher, Higher derivatives in type II and M-theory on Calabi-Yau threefolds. JHEP **1802**, 127 (2018). [arXiv:1702.08404 [hep-th]]
77. B.J. Broy, D. Ciupke, F.G. Pedro, A. Westphal, Starobinsky-type inflation from α'-corrections. JCAP **1601**, 001 (2016). [arXiv:1509.00024 [hep-th]]
78. M. Cicoli, D. Ciupke, S. de Alwis, F. Muia, α' Inflation: moduli stabilisation and observable tensors from higher derivatives. JHEP **1609**, 026 (2016). [arXiv:1607.01395 [hep-th]]
79. M. Berg, M. Haack, J.U. Kang, S. Sjors, Towards the one-loop Kahler metric of Calabi-Yau orientifolds. JHEP **1412**, 077 (2014). [arXiv:1407.0027 [hep-th]]
80. M. Haack, J.U. Kang, One-loop Einstein-Hilbert term in minimally supersymmetric type IIB orientifolds. JHEP **1602**, 160 (2016). [arXiv:1511.03957 [hep-th]]
81. M. Haack, J.U. Kang, Field redefinitions and Kahler potential in string theory at 1-loop. JHEP **1808**, 019 (2018). [arXiv:1805.00817 [hep-th]]
82. I. Antoniadis, Y. Chen, G. K. Leontaris, Perturbative moduli stabilisation in type IIB/F-theory framework. arXiv:1803.08941 [hep-th]
83. C.P. Burgess, A. Maharana, F. Quevedo, Uber-naturalness: unexpectedly light scalars from supersymmetric extra dimensions. JHEP **1105**, 010 (2011). [arXiv:1005.1199 [hep-th]]
84. J.P. Conlon, Moduli stabilisation and applications in IIB string theory. Fortsch. Phys. **55**, 287 (2007). [hep-th/0611039]
85. C.P. Burgess, C. Escoda, F. Quevedo, Nonrenormalization of flux superpotentials in string theory. JHEP **0606**, 044 (2006). [hep-th/0510213]
86. E. Witten, New issues in manifolds of SU(3) holonomy. Nucl. Phys. B **268**, 79 (1986)
87. C.P. Burgess, A. Font, F. Quevedo, Low-energy effective action for the superstring. Nucl. Phys. B **272**, 661 (1986)
88. M. Dine, N. Seiberg, Nonrenormalization theorems in superstring theory. Phys. Rev. Lett. **57**, 2625 (1986)
89. M. Cicoli, J.P. Conlon, A. Maharana, F. Quevedo, A note on the magnitude of the flux superpotential. JHEP **1401**, 027 (2014). [arXiv:1310.6694 [hep-th]]
90. M. Cicoli, F. Quevedo, R. Valandro, De Sitter from T-branes. JHEP **1603**, 141 (2016). [arXiv:1512.04558 [hep-th]]
91. A. Westphal, De Sitter string vacua from Kahler uplifting. JHEP **0703**, 102 (2007). [hep-th/0611332]
92. M. Cicoli, A. Maharana, F. Quevedo, C.P. Burgess, De Sitter string vacua from Dilaton-dependent non-perturbative effects. JHEP **1206**, 011 (2012). [arXiv:1203.1750 [hep-th]]
93. D. Gallego, M.C.D. Marsh, B. Vercnocke, T. Wrase, A new class of de Sitter Vacua in type IIB large volume compactifications. JHEP **1710**, 193 (2017). [arXiv:1707.01095 [hep-th]]
94. J. Biablack, D. Roest, I. Zavala, De Sitter Vacua from nonperturbative flux compactifications. Phys. Rev. D **90**(2), 024065 (2014). [arXiv:1312.5328 [hep-th]]
95. M. Cicoli, S. de Alwis, A. Westphal, Heterotic Moduli stabilisation. JHEP **1310**, 199 (2013). [arXiv:1304.1809 [hep-th]]
96. S.L. Parameswaran, S. Ramos-Sanchez, I. Zavala, On Moduli stabilisation and de Sitter Vacua in MSSM heterotic orbifolds. JHEP **1101**, 071 (2011). [arXiv:1009.3931 [hep-th]]
97. J.P. Conlon, A. Maharana, F. Quevedo, JHEP **05**, 109 (2009). https://doi.org/10.1088/1126-6708/2009/05/109 [arXiv:0810.5660 [hep-th]]
98. S. Krippendorf, M.J. Dolan, A. Maharana, F. Quevedo, JHEP **06**, 092 (2010). https://doi.org/10.1007/JHEP06(2010)092 [arXiv:1002.1790 [hep-th]]
99. M. Cicoli, C. Mayrhofer, R. Valandro, Moduli stabilisation for Chiral Global Models. JHEP **1202**, 062 (2012). [arXiv:1110.3333 [hep-th]]
100. M. Cicoli, S. Krippendorf, C. Mayrhofer, F. Quevedo, R. Valandro, D-Branes at del Pezzo singularities: global embedding and Moduli stabilisation. JHEP **1209**, 019 (2012). [arXiv:1206.5237 [hep-th]]
101. M. Cicoli, S. Krippendorf, C. Mayrhofer, F. Quevedo, R. Valandro, D3/D7 branes at singularities: constraints from global embedding and Moduli stabilisation. JHEP **1307**, 150 (2013). [arXiv:1304.0022 [hep-th]]
102. M. Cicoli, F. Muia, P. Shukla, Global embedding of fibre inflation models. JHEP **1611**, 182 (2016). [arXiv:1611.04612 [hep-th]]
103. M. Cicoli, I. Garcia-Etxebarria, C. Mayrhofer, F. Quevedo, P. Shukla, R. Valandro, Global orientifolded quivers with inflation. JHEP **1711**, 134 (2017). [arXiv:1706.06128 [hep-th]]
104. M. Cicoli, D. Ciupke, V.A. Diaz, V. Guidetti, F. Muia, P. Shukla, Chiral global embedding of fibre inflation models. JHEP **1711**, 207 (2017). [arXiv:1709.01518 [hep-th]]
105. L. Aparicio, M. Cicoli, S. Krippendorf, A. Maharana, F. Muia, F. Quevedo, JHEP **11**, 071 (2014). https://doi.org/10.1007/JHEP11(2014)071 [arXiv:1409.1931 [hep-th]]
106. L. Aparicio, M. Cicoli, B. Dutta, S. Krippendorf, A. Maharana, F. Muia, F. Quevedo, JHEP **05**, 098 (2015). https://doi.org/10.1007/JHEP05(2015)098 [arXiv:1502.05672 [hep-ph]]
107. A. Maloney, E. Silverstein, A. Strominger, De Sitter space in noncritical string theory. hep-th/0205316
108. E. Silverstein, Simple de Sitter solutions. Phys. Rev. D **77**, 106006 (2008). [arXiv:0712.1196 [hep-th]]
109. S. de Alwis, K. Givens, Physical Vacua in IIB compactifications with a single Kahler modulus. JHEP **1110**, 109 (2011). [arXiv:1106.0759 [hep-th]]
110. M. Rummel, A. Westphal, A sufficient condition for de Sitter vacua in type IIB string theory. JHEP **1201**, 020 (2012). [arXiv:1107.2115 [hep-th]]
111. B.S. Acharya, K. Bobkov, G. Kane, P. Kumar, D. Vaman, An M theory solution to the hierarchy problem. Phys. Rev. Lett. **97**, 191601 (2006). [hep-th/0606262]
112. B.S. Acharya, K. Bobkov, G.L. Kane, P. Kumar, J. Shao, Explaining the electroweak scale and stabilizing Moduli in M theory. Phys. Rev. D **76**, 126010 (2007). [hep-th/0701034]
113. C. Caviezel, P. Koerber, S. Kors, D. Lust, T. Wrase, M. Zagermann, On the cosmology of type IIA compactifi-

cations on SU(3)-structure manifolds. JHEP **0904**, 010 (2009). [arXiv:0812.3551 [hep-th]]
114. R. Flauger, S. Paban, D. Robbins, T. Wrase, Searching for slow-roll moduli inflation in massive type IIA supergravity with metric fluxes. Phys. Rev. D **79**, 086011 (2009). [arXiv:0812.3886 [hep-th]]
115. U.H. Danielsson, S.S. Haque, G. Shiu, T. Van Riet, Towards classical de Sitter solutions in string theory. JHEP **0909**, 114 (2009). [arXiv:0907.2041 [hep-th]]
116. C. Caviezel, T. Wrase, M. Zagermann, Moduli stabilization and cosmology of type IIB on SU(2)-structure orientifolds. JHEP **1004**, 011 (2010). [arXiv:0912.3287 [hep-th]]
117. U.H. Danielsson, P. Koerber, T. Van Riet, Universal de Sitter solutions at tree-level. JHEP **1005**, 090 (2010). [arXiv:1003.3590 [hep-th]]
118. U.H. Danielsson, S.S. Haque, P. Koerber, G. Shiu, T. Van Riet, T. Wrase, De Sitter hunting in a classical landscape. Fortsch. Phys. **59**, 897 (2011). [arXiv:1103.4858 [hep-th]]
119. D. Andriot, New constraints on classical de Sitter: flirting with the swampland. arXiv:1807.09698 [hep-th]
120. B.S. Acharya, F. Benini, R. Valandro, Fixing moduli in exact type IIA flux vacua. JHEP **0702**, 018 (2007). [hep-th/0607223]
121. J. McOrist, S. Sethi, M-theory and type IIA flux compactifications. JHEP **1212**, 122 (2012). [arXiv:1208.0261 [hep-th]]
122. B. de Carlos, A. Guarino, J.M. Moreno, Flux moduli stabilisation, supergravity algebras and no-go theorems. JHEP **1001**, 012 (2010). [arXiv:0907.5580 [hep-th]]
123. U. Danielsson, G. Dibitetto, On the distribution of stable de Sitter vacua. JHEP **1303**, 018 (2013). [arXiv:1212.4984 [hep-th]]
124. J. Blabäck, U. Danielsson, G. Dibitetto, Fully stable dS vacua from generalised fluxes. JHEP **1308**, 054 (2013). [arXiv:1301.7073 [hep-th]]
125. C. Damian, L.R. Diaz-Barron, O. Loaiza-Brito, M. Sabido, Slow-roll inflation in non-geometric flux compactification. JHEP **1306**, 109 (2013). [arXiv:1302.0529 [hep-th]]
126. C. Damian, O. Loaiza-Brito, More stable de Sitter vacua from S-dual nongeometric fluxes. Phys. Rev. D **88**(4), 046008 (2013). [arXiv:1304.0792 [hep-th]]
127. B.S. Acharya, A. Maharana, F. Muia, JHEP **03**, 048 (2019). https://doi.org/10.1007/JHEP03(2019)048 [arXiv:1811.10633 [hep-th]]
128. K. Choi, String or M-theory axion as a quintessence. Phys. Rev. D **62**, 043509 (2000). ([**hep-ph/9902292**])
129. N. Kaloper, L. Sorbo, Where in the String landscape is quintessence. Phys. Rev. D **79**, 043528 (2009). [arXiv:0810.5346 [hep-th]]
130. S. Panda, Y. Sumitomo, S.P. Trivedi, Axions as quintessence in string theory. Phys. Rev. D **83**, 083506 (2011). [arXiv:1011.5877 [hep-th]]
131. M. Cicoli, F.G. Pedro, G. Tasinato, Natural quintessence in string theory. JCAP **1207**, 044 (2012). [arXiv:1203.6655 [hep-th]]
132. J. Blabäck, U. Danielsson, G. Dibitetto, JCAP **03**, 003 (2014). https://doi.org/10.1088/1475-7516/2014/03/003 [arXiv:1310.8300 [hep-th]]
133. G. D'Amico, N. Kaloper, A. Lawrence, Phys. Rev. D **100**(10), 103504 (2019). https://doi.org/10.1103/PhysRevD.100.103504 [arXiv:1809.05109 [hep-th]]
134. Y. Olguin-Trejo, S. L. Parameswaran, G. Tasinato, I. Zavala, Runaway quintessence, out of the Swampland. arXiv:1810.08634 [hep-th]
135. M. Cicoli, G. Dibitetto, F.G. Pedro, JHEP **10**, 035 (2020). https://doi.org/10.1007/JHEP10(2020)035 [arXiv:2007.11011 [hep-th]]
136. A. Arvanitaki, S. Dimopoulos, S. Dubovsky, N. Kaloper, J. March-Russell, Phys. Rev. D **81**, 123530 (2010). https://doi.org/10.1103/PhysRevD.81.123530 [arXiv:0905.4720 [hep-th]]
137. S. Dimopoulos, S. Kachru, J. McGreevy, J.G. Wacker, JCAP **08**, 003 (2008). https://doi.org/10.1088/1475-7516/2008/08/003 [arXiv:hep-th/0507205 [hep-th]]
138. L. McAllister, E. Silverstein, A. Westphal, Phys. Rev. D **82**, 046003 (2010). https://doi.org/10.1103/PhysRevD.82.046003 [arXiv:0808.0706 [hep-th]]
139. Ruchika, K. Dutta, A. Mukherjee, A. A. Sen, Observational constraints on axion(s) with a cosmological constant. [arXiv:2005.08813 [astro-ph.CO]]