Four-dimensional de Sitter space is a Glauber-Sudarshan state in string theory

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ABSTRACT: We show that four-dimensional de Sitter space is a Glauber-Sudarshan state, i.e. a coherent state, over a supersymmetric solitonic background in full string theory. We argue that such a state is only realized in the presence of temporally varying degrees of freedom and after including quantum corrections, with supersymmetry being broken spontaneously. On the other hand, fluctuations over the resulting de Sitter space is governed by the Agarwal-Tara state, which is a graviton (and flux)-added coherent state. Once de Sitter space is realized as a coherent state, and *not* as a vacuum, its ability to remain out of the swampland as well as issues regarding its (meta)stability, vacuum energy, and finite entropy appear to have clear resolutions.

Contents

1 Introduction & setup

The search for four-dimensional de Sitter (dS) has led to claims varying from having numerous solutions [1–4] to having none [5–8]. The standard picture of a landscape of vacua [9], along with an anthropic principle to explain our own universe [10], runs into problems not because of any no-go theorem [11–14], or technical hindrances [15, 16], but due to obstructions against dS spacetimes ingrained in some fundamental aspects of quantum gravity [17–22]. These include trans-Planckian issues [23] challenging the very notion of a well-defined, unitary Wilsonian effective action for accelerating backgrounds [17–19], or that of instabilities associated with the vacuum state for dS leading to a *quantum swampland* [24]. It was recently shown that the hurdles related to an effective field theory description of dS, as summarized in [30], can, in principle, be overcome once *time-dependent* degrees are switched on [25–29]. In this work, we show that other fundamental obstructions are also avoided if four-dimensional dS space itself is regarded as a *coherent state* over a Minkowski spacetime. Although similar setups have been considered before [31–35], what is new here is an explicit construction of such a configuration realized in full string theory.[1] Remarkably, the resulting construction of dS space hints at natural resolutions to problems related to its stability, vacuum energy and entropy.

We shall focus on dS in the *flat* slicing, with a metric of the form:

$$ds^2 = \frac{1}{\Lambda |t|^2} \left(-dt^2 + dx_1^2 + dx_2^2 + dx_3^2 \right), \tag{1.1}$$

where Λ is the cosmological constant and the temporal coordinate t has a range $-\infty \le t \le 0$, with the late time regime given by $t \to 0$. We choose this specific realization of dS (1.1) only for computational efficiency and other realizations, specifically of Kasner-type, have

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[1]A more detailed version, with proofs and computations, appears in [36].

been considered earlier [25]. The question is how to realize a metric like (1.1) in, say, type IIB string theory? A natural first conjecture is the following metric configuration:

$$ds^2 = \frac{1}{\Lambda H^2(y)|t|^2}\left(-dt^2 + dx_i^2\right) + H^2(y)g_{MN}(y)dy^M dy^N, \tag{1.2}$$

where $H(y)$ is the warp-factor that only depends on the coordinate of the internal six-dimensional manifold whose *unwarped* metric is given by $g_{MN}(y)$. Crucially, note that the internal space is time-independent. To see what kind of fluxes are required to support a configuration like (1.2) in IIB string theory, we uplift the configuration to M-theory. This uplift is only for computational advantage, and has no deeper implications as the degrees of freedom remain unchanged. The uplifted metric becomes:

$$ds^2 = g_s^{-8/3}\eta_{\mu\nu}dx^\mu dx^\nu + g_s^{-2/3}H^2 g_{MN}dy^M dy^N + g_s^{4/3}|dz|^2, \tag{1.3}$$

where $g_s \equiv \sqrt{\Lambda}\,|t|\,H(y)$ is the IIA string coupling which is now a function of both y^M and t, implying that in M-theory the internal manifold becomes time-*dependent* and $g_s \to 0$ denotes *late time*. Additionally, since we require $(g_s/H) < 1$, our analysis would only make sense in the interval:

$$-\frac{1}{\sqrt{\Lambda}} < t < 0. \tag{1.4}$$

Beyond this limit, we lose all quantitative control on the dynamics. Consequently, this means that not only is our solution trustworthy only below the *quantum break-time* of dS [22] and the TCC limit [17], but it also provides a simple resolution of the coincidence problem of dark energy by limiting the age of our universe by the current Hubble time [37], and avoids problems of Boltzmann brains.

The dynamics is controlled by metric and fluxes because, to support such a manifold, we need G-fluxes on an eleven-dimensional space whose topology is:

$$\mathcal{M}_{11} = \mathbb{R}^{2,1} \times \mathcal{M}_8 \equiv \mathbb{R}^{2,1} \times \mathcal{M}_4 \times \mathcal{M}_2 \times \frac{\mathbb{T}^2}{\mathcal{G}}, \tag{1.5}$$

where (μ, ν) parametrize $\mathbb{R}^{2,1}$, (M, N) parametrize the six-dimensional base with unwarped metric $g_{MN}(y)$ and z^a parametrizes the toroidal fibre modded out by the isometry group $\mathcal{G}$. Let $\mathbf{G}_{MNPQ}$ and $\mathbf{G}_{MNab}$, where $(M, N) \in \mathcal{M}_6$ and $(a, b) \in \frac{\mathbb{T}^2}{\mathcal{G}}$, denote the 4-fluxes switched on in the internal space. Fluxes on a *compact* space cannot be arbitrary: they have to satisfy Gauss' law (or alternatively, cancel anomalies), solve the EOMs, and be *quantized*. The quantization condition is [38]:

$$\left[\frac{\mathbf{G}_4}{2\pi}\right] - \frac{p_1(y)}{4} \in \mathbb{H}^4\left(y, \mathbb{Z}\right), \tag{1.6}$$

where $\mathbf{G}_4$ is a four-*form* G-flux component, $p_1(y)$ is the first Pontryagin class and $\mathbb{H}^4(y, \mathbb{Z})$ is the fourth cohomology class. In a time-dependent internal space like (1.3), this quantization condition cannot be satisfied with time-*independent* G-flux components. Additionally, EOMs also require fluxes to become time-*dependent* raising, in turn, serious questions about how could these fluxes be quantized, and how could they satisfy Gauss' law. On the dual

IIB side, this means that the three and the five-form fluxes should become time-*dependent*, leading to similar questions. The axio-dilaton however remains time-independent because we are in the constant-coupling scenario of F-theory [25, 41]. In the absence of quantum corrections and time-dependent fluxes, as shown in [25], all the questions raised above cannot be answered and the system *is* inconsistent, stemming from the loss of g_s and M_p hierarchies, leading to a breakdown of effective field theory as another manifestation of the swampland conjectures [30]. More importantly, a metric like (1.3) is still a bit too constrained to be a solution, and what really works is a metric of the form:

$$
ds^2 = g_s^{-8/3}\eta_{\mu\nu}dx^\mu dx^\nu + g_s^{-2/3}H^2\Big(F_1(t)g_{\alpha\beta}dy^\alpha dy^\beta
$$
$$
+ F_2(t)g_{mn}dy^m dy^n\Big) + g_s^{4/3}|dz|^2, \tag{1.7}
$$

where $(\alpha, \beta) \in \mathcal{M}_2$ and $(m, n) \in \mathcal{M}_4$ of (1.5). The additional time-dependences governs how the six-dimensional manifold $\mathcal{M}_6 = \mathcal{M}_4 \times \mathcal{M}_2$ changes with respect to time. On the IIB side this converts (1.2) to the following:

$$
ds^2 = \frac{1}{\Lambda H^2(y)|t|^2} \left(-dt^2 + dx_1^2 + dx_2^2 + dx_3^2\right) \tag{1.8}
$$
$$
+ H^2(y)\Big(F_1(t)g_{\alpha\beta}(y)dy^\alpha dy^\beta + F_2(t)g_{mn}(y)dy^m dy^n\Big),
$$

making the internal space time-*dependent*. To preserve four-dimensional Newton's constant, we additionally require $F_1 F_2^2 = 1$ and both $F_i(t) \to 1$ as $g_s \to 0$. Imposing all these conditions, the system becomes tightly constrained but does have a solution, answering all the questions that we raised above, as shown in [25]. However, even after we get an effective potential which supports dS space in string theory, there still remains questions whether the radiative corrections, calculated for some vacuum state, leads to instabilities [24] or if the obstructions regarding trans-Planckian issues [17] lingers on. To answer these, we attempt to see if a metric like (1.8) or (1.7) can be realized as a *state* instead of a *vacuum*.

2 The Glauber-Sudarshan state

The state that we have in mind is the Glauber-Sudarshan state [39, 40], commonly called a coherent state, because such a state is closest to a *classical* configuration which can be realized in a quantum theory. Our aim would be to realize (1.7) as a Glauber-Sudarshan state over a supersymmetric solitonic vacuum. (Then, by dualization, (1.8) can also be realized as such a state.) The supersymmetric vacuum would be a warped Minkowski space. The supersymmetric solitonic background that we have in mind is of the form:

$$
ds^2 = \frac{1}{h^{2/3}(y)} \left(-dt^2 + dx_i^2\right) + h^{1/3}(y)g_{MN}^{(0)}dy^M dy^N, \tag{2.1}
$$

where $h(y)$ is the warp-factor, and can be supported by self-dual G-fluxes of the form $\mathbf{G}_{MNPQ}^{(0)}(y)$, as well as $\mathbf{G}_{0ijM}^{(0)}$, where $(M, N) \in \mathcal{M}_8$. Such a background has been studied in detail in [43]. If we study *fluctuations* over this background, they are classified by modes which may be easily determined. These modes will typically have non-trivial spatial

behavior, whose dynamics will be governed by a Schrödinger equation over a non-trivial potential, but their temporal behavior would be simple ($\sim e^{i\omega_{\mathbf{k}}t}$). This will help us to avoid any trans-Planckian issues, but new subtleties lie in the construction of the Glauber-Sudarshan state itself. In the original work of [39, 40], the state was created by shifting the *free* vacuum (or the harmonic vacuum) by a displacement operator. One of the main issue in our case is that there is *no* free vacuum in a highly interacting theory like M-theory! Secondly, M-theory has metric as well as G-flux components, so a Glauber-Sudarshan state would be more complicated to account for fluctuations of all these ingredients. Thirdly, even if we manage to construct such a state, how do we know that such a state survives the set of quantum corrections coming from perturbative, non-perturbative, non-local and topological interactions?

Let us start by answering the first question, related to the construction of the Glauber-Sudarshan state. Since there is no free vacuum once interactions are switched on, we only have the *interacting* vacuum $|\Omega\rangle$ to build our state from. We shall shift the interacting vacuum by a displacement operator and ask if this creates a state resembling the Glauber-Sudarshan state, namely:

$$|\sigma\rangle \equiv \mathbb{D}(\sigma)|\Omega\rangle, \tag{2.2}$$

at a specific time $t = t_0$. $\mathbb{D}(\sigma)$ is the *displacement* operator and $\sigma \equiv (\{\alpha_g\}, \{\beta_{\rm C}\})$ where $\{\alpha_g\}$ and $\{\beta_{\rm C}\}$ are the two sets of parameters associated with all the metric and the C-field components, respectively. The definition of a displacement operator in an interacting theory is, however, not clear. In the free theory, a displacement operator is constructed from annihilation $a_{\mathbf{k}}$ and creation $a_{\mathbf{k}}^\dagger$ operators for a given spatial momentum $\mathbf{k}$, so in an interacting theory we expect analogous operators $a_{\rm eff}(\mathbf{k})$ and $a_{\rm eff}^\dagger(\mathbf{k})$ to replace them, with $a_{\rm eff}(\mathbf{k})$ annihilating the interacting vacuum. Unfortunately this information is not enough to fix the form of $a_{\rm eff}(\mathbf{k})$, the latter being is a complicated function that mixes the free-field annihilation and creation operators for different spatial momenta [36]. Due to this, $\mathbb{D}(\sigma)$ is *not* a unitary operator anymore. However, there does exist one possible choice for $\mathbb{D}(\sigma)$ that not only fixes the form for $a_{\rm eff}(\mathbf{k})$, but also reproduces the background (1.7) as expectation values of the metric operators in the state (2.2). This choice works for any time t and is given by:

$$\mathbb{D}(\sigma, t) = \mathbb{D}_0(\sigma) \exp\left(i \int_{-T}^{t} d^{11}x\ \mathbf{H}_{\rm int}\right), \tag{2.3}$$

where $T \to \infty$ in a slightly imaginary direction; and $\mathbf{H}_{\rm int} \equiv \mathbf{H}_{\rm int}(g_{\rm MN}, {\rm C}_{\rm PQR})$ is the full interacting part of the M-theory Hamiltonian. $\mathbb{D}_0(\sigma)$ is the displacement operator for the harmonic vacuum, meaning that it displaces the harmonic vacuum by $\sigma \equiv \sigma(t)$ to create the required Glauber-Sudarshan state with one minor difference: it is the non-unitary part of the usual free vacuum displacement operator. In writing (2.3) we have ignored a multiplicative constant piece that is proportional to the *overlap* between the interacting and the harmonic vacuum, i.e. $\langle\Omega|0\rangle$. One can also work out the wavefunction in the configuration space for the state (2.2) satisfying (2.3). For example, say, for the spacetime

mode for the graviton, the wave-function of the state (2.2) may be expressed explicitly as:

$$\Psi_\Omega^{(\sigma)}(g_{\mu\nu}, t) = \exp\left[\int_{-\infty}^{+\infty} d^{10}\mathbf{k} \, \log\langle \widetilde{g}_{\mu\nu}(\mathbf{k}) | \Psi_{\mathbf{k}}^{(\sigma)}(t)\rangle\right], \tag{2.4}$$

where $|\Psi_{\mathbf{k}}^{(\sigma)}(t)\rangle$ is the Glauber-Sudarshan state in the Heisenberg representation for a given spatial mode $\mathbf{k}$ and $\widetilde{g}_{\mu\nu}(\mathbf{k})$ is the Fourier component of the graviton.

Indeed, this is all we need for the present purpose, as all of the background quantities simply appear by taking expectation values over the state (2.2) with the wave-function (2.4). For example, let us concentrate on the space-time metric again. The expectation value of the metric operator may be expressed as:

$$\langle \mathbf{g}_{\mu\nu} \rangle_\sigma = \frac{\int [\mathcal{D}g_{\mu\nu}] e^{i\mathbf{S}} \mathbb{D}^\dagger(\sigma) g_{\mu\nu} \mathbb{D}(\sigma)}{\int [\mathcal{D}g_{\mu\nu}] e^{i\mathbf{S}} \mathbb{D}^\dagger(\sigma) \mathbb{D}(\sigma)} = \frac{\eta_{\mu\nu}}{(\Lambda |t|^2 \mathrm{H}^2(y))^{4/3}}, \tag{2.5}$$

where $\mathbf{S} \equiv \mathbf{S}(g_{\mathrm{MN}}, \mathrm{C_{PQR}})$ is the total M-theory action. Since $\mathbb{D}(\sigma) \equiv \mathbb{D}(\sigma(t), t)$ is non-unitary, not only does it necessitate a division by another path integral as shown in the middle equality, but also it keeps the numerator from vanishing. In fact, $\mathbb{D}(\sigma)$ does what it is expected to do: it shifts the vacuum in such a way that the one-point functions do not vanish and ensures that they have the necessary expectation values. In addition, the choice (2.3) guarantees that there are no $\mathcal{O}\left(\frac{g_s^a}{\mathrm{M}_p^b}\right)$ corrections to (2.5) [36]. The above computation relies on two essential objects: (i) the space-time wave-functions (not the configuration space wave-functions!) which come from solving a class of Schrödinger equations with non-trivial potentials, and (ii) the Glauber-Sudarshan wave-function (2.4). Similarly, the expectation value of the G-flux components over the state (2.2) becomes:

$$\langle \mathbf{G}_{\mathrm{MNPQ}} \rangle_\sigma = \sum_p \mathcal{G}_{\mathrm{MNPQ}}^{(p)}(y) \left(\frac{g_s}{\mathrm{H}}\right)^{2p/3}, \tag{2.6}$$

where $(\mathrm{M, N}) \in \mathcal{M}_8$ with $p \in \frac{\mathbb{Z}}{2}$ and $p \geq \frac{3}{2}$. The latter condition stems from various criteria, including flux EOMs, Bianchi identities and subtleties with *localized fluxes*, as elaborated in [25]. We want to emphasize that the bound on p tells us that there are *no* time-independent fluxes allowed in this set-up.

Note that supersymmetry is broken spontaneously because the G-flux on the internal space is no longer self-dual, i.e.

$$|\langle \mathbf{G}_4 \rangle_\sigma - \langle *_8 \mathbf{G}_4 \rangle_\sigma| > 0, \tag{2.7}$$

where $\mathbf{G}_4$ is the four-form operator and $*_8$ is the Hodge dual with respect to the un-warped metric of the solitonic background (2.1). Note that the fluxes $\mathbf{G}_4^{(0)}$, supporting (2.1), are self dual and preserve supersymmetry for the solitonic vacuum (thereby canceling the zero-point energies) and, subsequently, supersymmetry is broken spontaneously by the Glauber-Sudarshan state. The cosmological constant is much lower than the energy scale of supersymmetry breaking due to suppression by a factor of the unwarped volume [36].

3 Fluctuations & the Agarwal-Tara state

Next, to study instabilities associated with QFT in dS, we wish to explore *fluctuations* over dS space in such a setting. Since dS itself is a state over the solitonic vacuum, the fluctuations should also come from a related state appearing as some deformation of the Glauber-Sudarshan state. It turns out that the required deformation is another well-known state, called the *Agarwal-Tara* state [45], or alternatively, as the graviton added coherent state. For us, this needs to be generalized in the same vein as the Glauber-Sudarshan state so that the Agarwal-Tara state will have to be both graviton and flux-added coherent state, which is given by:

$$\left| \Psi_{\mathbf{k}}^{(c_1 c_2)}(t) \right\rangle = \left[c_1 + c_2 \mathcal{G} \left(a_{\mathbf{k}} + a_{\mathbf{k}}^{\dagger}; t \right) \right] \left| \Psi_{\mathbf{k}}^{(\sigma)}(t) \right\rangle, \tag{3.1}$$

where we have used (2.3), and denote $a_{\text{eff}}(\mathbf{k})$ and $a_{\text{eff}}^{\dagger}(\mathbf{k})$ by $a_{\mathbf{k}}$ and $a_{\mathbf{k}}^{\dagger}$ to simplify the notation, and c_i are constants with $|c_2| \ll |c_1|$. We have further restricted to the graviton sector to avoid over-burdening the formula with complications from the flux sector. $\mathcal{G}(w, t)$ could be thought of as a polynomial function of w with time-dependent coefficient (details appear in [36]). Subtleties aside, what is significant for us is that the expectation of the graviton operator in such a state gives us:

$$\langle \mathbf{g}_{\mu\nu} \rangle_{\Psi^{(c_1 c_2)}} = \frac{\eta_{\mu\nu}}{(\Lambda |t|^2 \mathrm{H}^2(y))^{4/3}} \tag{3.2}$$
$$+ c_2 \int d^{11}k \; h(\mathbf{k}, k_0) \; \alpha_{\mu\nu}(\mathbf{k}, k_0) \; \psi_{\mathbf{k}}(\mathbf{x}, y, z) \; e^{ik_0 t},$$

where $\alpha_{\mu\nu}$ is precisely the parameter from the set $\{\alpha_g\}$ of $\sigma \equiv (\{\alpha_g\}, \{\beta_C\})$ that defines the Glauber-Sudarshan state (2.2). $\psi_{\mathbf{k}}(\mathbf{x}, y, z)$ is the eleven-dimensional spatial wave-function that appears from the Schrödinger equation over the solitonic background alluded to earlier. We have also taken $c_1 = 1$, so only c_2 appears in (3.2). $h(\mathbf{k}, k_0)$ is the Fourier component that comes from $\mathcal{G}(w, t)$ function in (3.1). The crucial take-home point from (3.2) is that the temporally varying frequencies $\omega_{\mathbf{k}}(t)$ that we get from fluctuations over a dS *vacuum* are nothing but *artifacts* of the Fourier transforms over the Glauber-Sudarshan state. In other words:

$$\exp\left(-i\omega_{\mathbf{k}}(t) t \right) \equiv \int dk_0 \; h(\mathbf{k}, k_0) \; e^{-ik_0 t}, \tag{3.3}$$

where $\omega_{\mathbf{k}}(t)$ is in general a complex function. Such a conclusion seems to point out that there are no trans-Planckian censorship required for our construction [17] because all modes originate secretly from fluctuations over our solitonic background (2.1). (We also satisfy the conditions of having a underlying Lorentz-invariant spacetime, with a local vacuum, required to avoid the original trans-Planckian problem [23].) As a result, Wilsonian analysis *can* be performed for these fluctuations because all modes in our theory can be expressed as linear combinations of the ones over the solitonic background with time-*independent* frequencies. Since our dS space, itself, is a state which results from these modes, it is no surprise that we also get fluctuations on top of dS as another consequence of them. However, we still need to demonstrate the stability of our construction, which we shall address next.

4 Stability

How do we know that the quantum corrections do not take us away from the background (1.7) and G-flux (2.6) configurations? In fact, any small changes to the $F_i(t)$ factors in (1.7) will switch on some time-dependence of the Newton's constant. In order to avoid this, we have to establish the stability of the solution. This also gets tied up with the equations of motion in the presence of *all* possible quantum corrections. It turns out that the background equations may be presented most succinctly as Schwinger-Dyson equations (SDEs). The SDEs [46–48] are expressed as expectation values over the Glauber-Sudarshan state and therefore suits our construction very well since we have considered the metric, the G-fluxes as well as the fluctuations as expectation values. The SDEs for our case may be divided into two set. One set is easy to write down and is given by:

$$\frac{\delta \mathbf{S}^{(\sigma)}}{\delta \langle g^{\mathrm{MN}} \rangle_\sigma} = \frac{\delta \mathbf{S}^{(\sigma)}}{\delta \langle \mathrm{C}^{\mathrm{MNP}} \rangle_\sigma} = 0, \tag{4.1}$$

where $\mathbf{S}^{(\sigma)} \equiv \mathbf{S}^{(\sigma)} \left(\langle g_{\mathrm{MN}} \rangle_\sigma, \langle \mathrm{C}_{\mathrm{PQR}} \rangle_\sigma \right)$ differs from $\mathbf{S} = \mathbf{S}(g_{\mathrm{MN}}, \mathrm{C}_{\mathrm{PQR}})$ in

$$\langle \mathbf{g}_{\mu\nu} \rangle_\sigma = \frac{\int [\mathcal{D} g_{\mu\nu}] e^{i\mathbf{S}} \mathbb{D}^\dagger(\sigma) g_{\mu\nu} \mathbb{D}(\sigma)}{\int [\mathcal{D} g_{\mu\nu}] e^{i\mathbf{S}} \mathbb{D}^\dagger(\sigma) \mathbb{D}(\sigma)} = \frac{\eta_{\mu\nu}}{(\Lambda |t|^2 \mathrm{H}^2(y))^{4/3}}, \tag{4.2}$$

by the appearance of the expectation values in the integrands themselves. The second set of equations are however more involved and they include both the Faddeev-Popov ghosts and the displacement operator [36]. For us (4.1) will suffice, as the form of these equations precisely imply the equations of motion already studied in [25] and, therefore, pursues the path laid down there, as follows. Once we express the equations order by order in $\left(\frac{g_s}{\mathrm{H}} \right)$, the zeroth order equations precisely determine the background (1.7) with the G-fluxes as in (2.6). Going to the higher orders in $\left(\frac{g_s}{\mathrm{H}} \right)$ then switches on three things: (a) higher orders, i.e $p > \frac{3}{2}$ in (2.6) for the G-fluxes, (b) higher order terms for the $F_i(t)$ factors in (1.7), and (c) higher order quantum terms discussed in [25]. Together, they *balance* each other in such a way that the zeroth order metric and G-fluxes, from (1.7) and (2.6) respectively, *do not receive any corrections* [36].

We can also see how moduli can be stabilized in such time-dependent background. The time-independent fluxes and the quantum corrections that are required to support the solitonic background (2.1) also stabilize the Kähler and the complex structure moduli of the internal manifold *without* breaking any supersymmetry. The time-dependent background can then be generated by taking the expectation values of the metric and the G-flux operators as in (4.2) and (2.6). This means the stabilized moduli at the solitonic level now vary in a controlled way as the internal metric (1.7) evolves in time. Thus at every instant of time the moduli are stabilized with no Dine-Seiberg runaway [50]. We call this *dynamical moduli stabilization*.

An advantage of choosing M-theory, as against type IIB theory, to do our computations is that we are, in fact, taking into account the full string spectra, which includes the NS and RR sectors of IIB. In M-theory the graviton state, when dimensionally reduced over the 11-dimensional cycle, reproduces the zero mode gravitons. On the other hand, most

of the massive gravitons coming out of stringy excitations decay at the tree and one-loop level so that they do not contribute to the dynamics of the coherent states. The remaining has been integrated out and they contribute to the infinite tower of perturbative and non-perturbative, including non-local and topological, quantum corrections. Surprisingly, it has been shown that the coherent state is *stable* under all these contributions from higher order curvature (and flux) terms that come out of the combination of the zero mode gravitons, the instanton gases and the perturbative loops [36]. The action $\mathbf{S}^{(\sigma)}$ above contains all types of quantum corrections, for instance, non-perturbative corrections of the BBS [49] and KKLT [1] type instantons as well as the action of the branes and surfaces, including fermionic interactions, as shown in [36]. Taking this exhaustive array of quantum corrections allowed us to venture to the regime where the curvature could be more than the string scale. In such a scenario, a countable (but large) number of perturbative and nonperturbative quantum corrections become essential [36]. Indeed, this is an important point since the cosmological constant in four-dimensions, in our formalism, emerges from a delicate balancing between the fluxes and the quantum corrections but *without* any vacuum energies [25]. Furthermore, since the fluctuations are ultimately built out of the interacting vacuum of the solitonic background, perturbative quantum corrections do not lead to any instabilities as everything is expressed in terms of time-independent mode expansions [36], as opposed to the case of classical dS [24, 30]. Although our analyses demonstrates the stability of our state against quantum corrections, we have not shown if the conditions necessary for constructing this Glauber-Sudarshan state are generic or not. Put differently, we have not examined whether this state is an attractor on the space of all such states constructed on top of the warped-Minkowski background, which we leave for future work.

As already mentioned, what makes our solution more physically interesting is that it resolves instabilities related to QFT of fluctuations on top of classical dS spacetime [24], which are not unique to stringy constructions. Remarkably, these fluctuations also turn out to be intimately connected to the *entropy* of dS spacetime, an old puzzle being the origin of its finiteness within the context of full quantum gravity [9]. Specifically, in our solution, a natural resolution would be in interpreting this as the *entanglement entropy* between the fluctuation modes, on top of the time-independent solitonic vacuum, which give rise to the Glauber-Sudarshan state itself. The way we get a finite von Neumann entanglement entropy, due to this mode-coupling, is that we have a reduced density matrix corresponding to tracing out the modes which are super-horizon, and treating the causal patch of an inertial observer as our system. The key observation is that if the interactions could be turned off, such an entanglement entropy would consist of diverging parts alone but, in that case, we would have no dS space either as $\mathbf{H}_{\text{int}}$ is crucial in the construction of (2.4). It had been heuristically argued earlier that, for a coherent state description of dS, a finite number of gravitons is synonymous to the finite entropy associated with it; furthermore, it was also noted that interactions between the gravitons are important to arrive at this conclusion [22]. In our case, we sketch out the concrete reason as to how one arrives at a finite number of highly-interacting gravitons (and flux particles) for our Glauber-Sudarshan state necessary for having a finite entropy for the resulting dS space. Our solution also satisfies the expectation of the dS symmetries being emergent for

a 'reasonably' short time-period, and are *not* eternal, which is a necessary condition for having a finite entropy [52, 53]. Without going into details, note that the above procedure is technically tractable due to the fact that the entanglement entropy corresponding to coherent states coincide with that for the vacuum [36, 54].

5 Conclusion

To conclude, in this work, we have shown how the EOMs can be solved in full string theory to obtain a dS spacetime by avoiding the swampland, once time-dependent degrees of freedom are turned on, and including quantum corrections, by constructing the solution in terms of a Glauber-Sudarshan state. Remarkably, this state (a) solves the trans-Planckian censorship problem, (b) is stable against both perturbative and non-perturbative quantum corrections, and (c) provides a microscopic understanding of its entropy as an inherently quantum quantity. Furthermore, the order of magnitude of the time-limit before which the system becomes strongly-coupled gives an easy way out of the coincidence problem of cosmology as well as an escape from issues related to Boltzmann brains. This much shorter timescale is a distinguishing feature of our model as opposed to the lifetime of dS in compactifications such as KKLT scenario [1].

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