Swampland Conjectures for an Almost Topological Gravity Theory

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We analyze AdS and dS swampland conjectures in a three-dimensional higher spin theory with self-interacting matter, which contains conformal gravity and is almost topological. A theory of a similar type was proposed as the effective theory in the high energy phase of non-critical M-theory in 3D. With some details differing from the usual string theory story, it is found that the resulting effective theory, namely topologically massive gravity, fits well into the web of the proposed swampland conjectures. Supporting a recent proposal, in particular we find that this 3D theory gives rise to a quantum break time that scales like \( t_Q \sim H^{-1} \).

I. INTRODUCTION

The swampland program [1] implies a paradigm shift in the field of string phenomenology. Instead of analyzing the string theory landscape for certain well understood classes of string compactifications, one tries to mark the boundary between the landscape and the swampland. This is quantified in the form of a number of swampland conjectures, which are supported by concrete string models and/or by more general quantum gravity arguments (see [2–4] for recent reviews). In this endeavor, our intuition is led by experience from critical string theory or M-theory and their low-energy effective actions. As such, we have a universal gravity sector that is governed by an Einstein-Hilbert action that upon quantization leads to gravitons interacting with a strength set by the Planck mass.

The purpose of this letter is to consider effective gravity theories that are of a different kind but have still appeared in the string theory literature, hence having a good chance to be in the string landscape. We want to challenge such theories with some of the swampland conjectures, in particular the AdS and dS swampland conjectures, as well as the proposal of quantum breaking.

One instance where such gravity theories of a different kind are expected to appear is in the high energy or high temperature regime of string theory. Unfortunately, for temperatures beyond the Hagedorn transition these theories are not completely understood, but there are indications, like the quadratic scaling of the free-energy with temperature [5], that the number of degrees of freedom gets reduced. This letter is also motivated by the recent proposal [6] that in the high temperature regime of string theory, the quantum break time should scale like \( t_Q \sim H^{-1} \) irrespective of the number of dimensions. Recall that, as shown in [7], for Einstein gravity in d space-time dimensions this scaling is \( t_Q \sim M_{pl}^{d-2}/H^{d-1} \).

Clearly, for making progress we need a concrete and well treatable starting ground. For our purposes, this is provided by the non-critical M-theory in 3D proposed by Hořava/Keeler [8–10] which is exactly solvable in terms of a non-relativistic Fermi liquid in 2 +1 dimensions. This is motivated by the known tachyon condensation from ten-dimensional type 0 theories, which represent the infinite temperature limit of type II theories, to the two-dimensional type 0 theories [11]. The non-critical M-theory arises from the thermal version of this two-dimensional theory when D0-branes are included [8–10]. In the conformal limit this theory was supposed to be described by an effective higher spin generalization of conformal gravity equipped with a higher spin matter field. The ground state of non-critical M-theory corresponds to an AdS\(_3 \times S^1\) solution of this theory plus a massless fermion. Without the matter field, the theory would be topological with no propagating degrees of freedom. Note that such topological theories were also invoked in the recent proposal [12] of a topological phase in the early universe.

This higher spin theory describing the conformal limit of non-critical M-theory is motivating our choice of a non-standard effective gravity theory. The latter is obtained as a truncation to just conformal gravity, where we will also include a conformally coupled, scalar matter field\(^1\). This is expected to decouple in the strict conformal limit, but is still present at large but finite energies. By also including a conformal self-interaction of the scalar field, this truncated theory admits non-trivial AdS\(_3\) and dS\(_3\) solutions for a non-vanishing vacuum expectation value of the scalar. The resulting effective theory is known as topologically massive gravity (TMG) [13], carrying the wrong sign of the Einstein-Hilbert term. In this letter we consider this almost topological theory as a toy model for the high energy phase of string theory and analyze whether it does still satisfy in particular those swampland conjectures dealing with dS and AdS solutions.

\(^1\) Note that the truncation may in principle have qualitatively different properties from its higher spin completion. In particular, the much bigger gauge invariance of the higher spin theory may render some solutions pure gauge. We thank the referee for raising this point.
II. NON-CRITICAL M-THEORY IN 3D

Three-dimensional non-critical M-theory was proposed in [8, 9] as the strong coupling limit of non-critical type 0A theory in 2D. In [8] a non-perturbative definition of this theory was formulated as a double-scaled non-relativistic Fermi liquid in 2 + 1 dimensions. This theory is closely related to the matrix model description of 2D type 0A theory by summing over all D0-brane charges. Moreover, it has an infinite dimensional symmetry described by a bosonic higher-spin Lie algebra, called \( W_0 \) in [9]. Important in the following is that the conformal algebra \( SO(3, 2) \) is a finite dimensional subalgebra of \( W_0 \). For more details we refer the reader to the original literature.

In [9] it was found that in the high-temperature regime the non-perturbative partition function is closely related to the one of the topological A-model on the resolved conifold. In the conformal limit \( (\alpha' \rightarrow \infty) \) it was argued in [10] that the ground state of non-critical M-theory should be an \( AdS_2 \times S^1 \) space-time accompanied by a massless free fermion as the only propagating degree of freedom.

As proposed in [10], a higher spin theory admitting such a solution is given by the higher spin Chern-Simons theory [14]

\[
S_{HCS} = \frac{1}{4} \int_M \text{Tr} \left( A \wedge dA + \frac{2}{3} A \wedge A \wedge A \right) \quad (1)
\]

with the gauge field one-form \( A \) taking values in the infinite-dimensional higher-spin Lie algebra \( W_0 \). As shown in [15], matter fields can be coupled to this topological theory at least at the level of the matter field equations of motion in a CS background field \( A \) via

\[
d(|\Phi \rangle + A \star |\Phi \rangle) = 0, \quad (2)
\]

where \( |\Phi \rangle \) is a higher spin matter field that contains a scalar and a fermion as the lowest components. In an \( AdS_2 \times S^1 \) (or also \( AdS_3 \) ) background with all higher spin fields set to zero, all these equations reduce to two equations of motion. One is the Klein-Gordon equation for a conformally coupled massless scalar field and the other is the Dirac equation for a massless fermion. Because of the absence of a propagating scalar degree of freedom in the conformal limit, the scalar field was set to zero by hand in [10]. The scalar is related to the “tachyon” in the effective 2D type 0A theory and is only expected to completely decouple in the strict \( \alpha' \rightarrow \infty \) limit, which could be made manifest by setting \( \phi = \varphi/(\alpha')^{1/4} \) in the following.

A. Conformal gravity with matter

The theory of interest in this paper is a truncation of this higher spin theory where we consider just the spin-2 mode and the scalar and fermionic matter fields. Actually, the fermion will not play any role in the following. Recall that restricting the higher spin CS theory (1) to just the spin-2 mode one gets conformal gravity in 2+1 dimensions [16], a topological gravity theory governed by the Chern-Simons action

\[
S_{CS}^{SO(3,2)} = \frac{1}{4} \int_M \text{Tr} \left( A \wedge dA + \frac{2}{3} A \wedge A \wedge A \right) \quad (3)
\]

with the one-form \( A \) taking values in the 3D conformal algebra \( SO(3, 2) \). The resulting equation of motion says that the Cotton tensor is vanishing

\[
C_{\mu \nu} = \epsilon_{\mu}^{\alpha \beta} \nabla_{\alpha} \left( R_{\beta \nu} - \frac{1}{4} g_{\beta \nu} R \right) = 0, \quad (4)
\]

which is satisfied for a conformally flat metric. Note that being topological, this theory does not admit propagating degrees of freedom. Moreover, by being conformal it does not have a dimensionful gravitational coupling like the Planck mass.

Now, one can couple this gravity theory to fermionic and bosonic matter fields. Here we focus on a scalar field, whose action takes the following form

\[
S_{\text{eff}} = S_{CS}^{SO(3,2)} + S_{\text{ferm}} - \int d^3 x \sqrt{-g} \left( g_{\mu \nu} \partial_\mu \phi \partial_\nu \phi \right) + \frac{\xi_3}{2} R \phi^2 + \frac{\kappa}{6} \phi^6 \quad (5)
\]

with \( \xi_3 = 1/8 \) and where we also added a conformal, sixth order self-coupling of the field \( \phi \) with a dimensionless coupling constant \( \kappa \). From now on we suppress the contribution from the fermionic matter field. One can show that this action is indeed invariant under a local Weyl-transformation

\[
g_{\mu \nu} \rightarrow e^{2\Lambda(x)} g_{\mu \nu}, \quad \phi \rightarrow e^{-\Lambda(x)/2} \phi. \quad (6)
\]

It was argued in [17] that also for the higher spin generalization, such an interaction term can be consistently introduced. Varying now \( S_{\text{eff}} \) with respect to the metric, one arrives at the gravity equation of motion \( C_{\mu \nu} - T_{\mu \nu}^{\text{bos}} = 0 \), that in detail takes the form

\[
C_{\mu \nu} - \frac{\phi^2}{8} \left( R_{\mu \nu} - \frac{1}{2} g_{\mu \nu} R \right) + \frac{\kappa}{6} \phi^6 g_{\mu \nu} - \left( \partial_\mu \phi \partial_\nu \phi - \frac{1}{2} g_{\mu \nu} (\partial \phi)^2 \right) - \frac{1}{4} \left( \phi \Box \phi + (\partial \phi)^2 \right) g_{\mu \nu} \quad (7)
\]

\[
+ \frac{1}{4} \left( \phi \nabla_\mu \nabla_\nu \phi + \partial_\mu \phi \partial_\nu \phi \right) = 0.
\]

It is remarkable that the Einstein tensor reappears in the equation of motion, though not from varying the Chern-Simons gravity action but from varying the action for the conformally coupled scalar. However, the sign in front of the Einstein tensor is opposite so that effectively gravity acts repulsively. If \( \phi \) carries a non-vanishing vacuum expectation value \( \phi_0 \), then the prefactor of the Einstein tensor can be thought of as an effective Planck scale

\[
\tilde{M}_{pl} = -\frac{\phi_0^2}{8}. \quad (8)
\]
Taking the trace of (7) one finds an expression that vanishes if the equation of motion for the scalar field is obeyed
\[ \Box \phi - \frac{1}{8} R \phi - \kappa \phi^5 = 0. \] (9)

This is of course a consequence of the underlying Weyl symmetry.

B. AdS and dS solutions

Let us consider potential AdS$_3$ and dS$_3$ solutions of radius $\ell$. These are conformally flat and have a constant Ricci scalar, so the equation of motion for $\phi$ admits the non-trivial solution
\[ \phi_0^4 = \frac{3}{4|\kappa| \ell^2}. \] (10)

Here AdS$_3$/dS$_3$ has curvature $R = \mp 6/\ell^2$ so that one needs $\kappa > 0$ for AdS$_3$ and $\kappa < 0$ for dS$_3$. One can show that also the gravity equation of motion (7) is satisfied.

Since $\ell$ sets a length scale, this can be thought of as breaking the Weyl symmetry spontaneously, with the flat direction in the ($\phi, \ell$) plane giving a Goldstone mode. Note that conformal gravity equipped with an additional scalar field with non-vanishing vacuum expectation value $\phi_0 \neq 0$ is nothing else than the well known theory of topologically massive gravity (TMG), whose action is usually written as
\[ S_{TMG} = \frac{\widehat{M}_{pl}}{2} \int d^3x \left[ \sqrt{-g} (R - 2\Lambda) + \frac{1}{2\mu} \epsilon^{\mu\nu\rho} \left( \Gamma^\gamma_{\mu\beta} \partial_\nu \Gamma^\beta_{\rho\alpha} + \frac{2}{3} \Gamma_\gamma^{\alpha\beta} \Gamma^\gamma_{\rho\alpha} \Gamma^\beta_{\nu\beta} \right) \right]. \] (11)

Note that the Weyl symmetry can be used to eliminate fluctuations of the field $\phi$. By comparison, one finds a negative Planck scale $\widehat{M}_{pl} = M_{pl}$ and
\[ \mu = \frac{\phi_0^2}{8} = \frac{1}{16\ell} \frac{3}{|\kappa|}, \quad \Lambda = \mp \frac{1}{\ell^2}. \] (12)

This theory was first introduced in [13], where it was shown that around a Minkowski solution it has one massive propagating spin-2 degree of freedom of mass $m_{(2)} = \mu$. Thus, as expected for a local symmetry, the Goldstone mode does not survive as a massless mode.

III. SWAMPLAND CONJECTURES

In this section we confront this effective almost topological theory with various swampland conjectures, in particular those dealing with AdS and dS solutions.

A. AdS swampland conjectures

Thus, we consider the case $\kappa > 0$ so that the action allows us to read off an effective potential
\[ V(\phi) = -\frac{3}{8\ell^2} \phi^2 + \kappa \phi^6. \] (13)

The form of this potential is shown in figure 1, where the non-trivial solution (10) for $\phi_0$ is evident.

![FIG. 1. The potential $V(\phi)$ for $\kappa = 1$ and AdS length scale $\ell = 10$.](image)

It can readily be checked that the AdS minimum satisfies the AdS moduli scalar separation conjecture [18] which in our case says $m_0 \ell < c$, for some $O(1)$ constant $c$. Just from $V(\phi)$ we find $m_0^2 = 3/\ell^2$.

Next, we take a look at the AdS distance conjecture [19]. From (12) we have for the cosmological constant $|\Lambda_{AdS}| \sim \phi_0^4$. Therefore it vanishes for $\phi_0 = 0$ which, opposed to the AdS distance conjecture, is at finite distance in the moduli space.

The AdS distance conjecture also says that there should be a tower of light states with scaling $m \sim |\Lambda|^{\alpha}$, $\alpha > 0$. From the (refined) swampland distance conjecture [20, 21], one expects a tower of states with masses
\[ m = m_0 e^{-\Delta_\phi/\sqrt{|\Lambda_{pl}|}} \sim \phi_0^2 \sim \frac{1}{\ell} \sim \Lambda_{AdS}^{1/4} \] (14)

where we have used that the only mass-scale in the problem is $m_0 \sim \phi_0^2$ and the effective Planck scale (8).

Since non-critical M-theory is defined in 3D, there is no compactification and hence no tower of Kaluza-Klein modes involved here. A natural candidate for this tower of states are the infinitely many higher spin fields in the Chern-Simons action. For the spin-2 field, the mass of the single massive degree of freedom of TMG with a negative cosmological constant was computed in [22, 23]
\[ m_{(2)}^2 = \left( \mu + \frac{2}{\ell} \right)^2 - \frac{1}{\ell^2} = \frac{1}{\ell^2} (\mu \ell + 3)(\mu \ell + 1) \] (15)

with $\mu \ell = \sqrt{3/(256\kappa)}$. This can be thought of as the $\mu$-corrected mass of the field $\phi$ that we read off from the effective scalar potential (13) as $m_0^2 = 3/\ell^2$. Moreover, the relation (15) correctly reproduces the mass of the graviton in the Minkowski limit $\ell \to \infty$. 


Since we have $\mu \ell > 0$, this mass can never vanish. Thus, our case is different from the chiral model discussed in [24], which in contrast to our case had a positive Planck mass. Taking some results for higher spin topologically massive gravity [25–27] into account, a natural conjecture for the mass of the higher spin fields would be

$$m^2_{(s)} = \frac{(s-1)}{\ell^2} \left( (s-1)\mu \ell + (s+1) \right) \left( \mu \ell + 1 \right), \quad (16)$$

which for $s \gg 1$ indeed gives the desired scaling $m_{(s)} \sim s \mu \sim s \phi_0^6$. Of course, a more thorough computation in the framework of the full $W_0$ higher spin theory is necessary to really confirm this relation.

Therefore, consistent with the finite distance in moduli space, there is only a polynomially and not an exponentially light tower of states. This is analogous to a similar polynomial scaling found for the finite distance conifold point in [28]. Similar to that case, applying the emergence proposal [29–31] will lead us to a field dependent UV cut-off

$$\Lambda_{UV} \sim \phi_0^5 \sim \frac{1}{\ell} \quad (17)$$

so that the number of light species is constant and does not depend on the field $\phi$. This cut-off is also consistent with the spin-2 swampland conjecture proposed in [32].

For 3D gravity with negative cosmological constant, it is known that there exist BTZ black-holes which are locally isometric to $AdS_3$ space. Therefore, they will also be solutions of TMG and also of the full equations of motion (7) and (9). However, with the wrong sign of the Einstein-Hilbert term they will carry negative energy and therefore signal a non-perturbative instability of the $AdS_3$ background. Let us remark that this is consistent with the proposed instability [33] of non-supersymmetric $AdS$-backgrounds.

To summarize, up to the finite distance in moduli space, this almost topological higher spin theory has all the ingredients to satisfy the $AdS$ swampland conjectures.

### B. $dS$ swampland conjecture

Now let us choose $\kappa < 0$ and consider the non-trivial $dS_3$ solution. Writing $\ell = 1/H$ the action features a potential

$$V(\phi) = \frac{3}{8} \frac{H^2 \phi^2}{|\kappa|} - \frac{1}{6} \phi^6 \quad (18)$$

which is just the one from figure 1 flipped at the $x$-axis. Therefore, we have a $dS$ maximum and can check whether the (refined) $dS$ swampland conjecture [34–36] in 3D is satisfied, i.e. in particular the relation

$$M_{pl}|V''| > c' V . \quad (19)$$

Computing both sides we find $V(\phi_0) = \frac{1}{4} \phi_0^6$ and $V''(\phi_0) = -4\kappa \phi_0^4 = -3/\ell^2$. Using now the effective Planck scale (8) we get that both sides scale as $\phi_0^6$ and that (19) is satisfied if $c' < 3/2$. The value of $c'$ is not known. One proposal [6] is $c' = 1/(D-1)(D-2)$ which for $D = 3$ yields $c' = 1/2$.

### C. Quantum break time

It has been proposed in [7, 37, 38] that $dS$ solutions in quantum gravity experience the phenomenon of quantum breaking. This is argued to come from the decoherence of the $dS$ bound state of graviton species.

First we notice that since $dS_3$ is conformally flat, it is a solution to pure conformal gravity and without any additional propagating modes present it would be eternal and no quantum breaking would occur. However, after conformally coupling it to a scalar field, one finds a massive spin-2 degree of freedom so that a priori quantum breaking can occur.

Since the only scale in the problem is set by $H$ or $\phi_0^6$, respectively, just by dimensional reasoning the quantum break time should scale as $t_Q \sim H^{-1}$. At first sight this is different from the quantum break time $t_Q \sim M_{pl}/H^2$ obtained in [7] for a three-dimensional Einstein gravity theory. However, taking into account the effective Planck scale (8) of TMG, the two expressions indeed agree.

This finding is consistent with the proposal in [6], namely that in the high temperature regime of string theory, the quantum break time should scale like $t_Q \sim H^{-1}$ irrespective of the number of dimensions. Due to the censorship of quantum breaking, there should be a classical mechanism leading to a faster decay. From the form of the potential in figure 1 it is clear that there exists a tachyonic instability (in the $\phi$ direction) of mass $m^2 = -3H^2$ leading indeed to a decay time $t_{dec} \sim H/|m|^2 \sim H^{-1}$.

### D. Higuchi bound

Finally, let us make a short comment on the Higuchi bound [39, 40]. In particular, it says that in a $d$-dimensional $dS$ background there exists a lower bound for the mass of an helicity $t$ mode of a massive field of spin $s$

$$m^2_{(s,t)} \geq H^2(s-t-1)(s+t+d-4) . \quad (20)$$

Since in our case the sign of the Einstein-Hilbert term is reversed, we expect that ghost freedom in $dS$ is guaranteed in the opposite regime, i.e. for the helicity $t = s$ mode in 3D we get

$$m^2_{(s,s)} \leq -H^2(2s-1) . \quad (21)$$

Next we need to know the mass of the spin-2 mode for TMG. Redoing the computation from [22, 23] for positive
cosmological constant $\Lambda = \ell^{-2} = H^2$, we find
\[
m^2_{(2)} = -\left(\mu + \frac{2}{\ell}\right)^2 + \frac{1}{\ell^2},
\]
which correctly reproduces the mass of the tachyonic mode in the $\mu \ell \to 0$ limit. This calls for the higher spin generalization
\[
m^2_s = -\left(\left(s - 1\right)\mu + \frac{s}{\ell}\right)^2 + \frac{1}{\ell^2} \leq (-s^2 + 1)H^2.
\]
which for $s \geq 2$ indeed satisfies the (sign reversed) Higuchi bound (21).

IV. CONCLUSIONS

In this letter we have studied a well motivated (toy) model for the high energy phase of string theory regarding its consistency with some of the swampland conjectures. We considered the truncation of the higher spin Chern-Simons theory to just the spin-2 mode equipped with a conformal, self-interacting scalar field. For this model we found $AdS$ and $dS$ vacua that, up to some slight modifications, have a good chance to satisfy the known swampland criteria. Some steps like the determination of the masses of the higher spin modes were left open. It would also be interesting to consider similar models of this type, as for instance the four-dimensional Weyl-squared action [41–43].

The result of this letter can be interpreted in two ways. First, it demonstrates that there is a good chance for the existence of non-standard effective gravity theories in the landscape of string theory. Second, it gives credence to the proposal that the high energy phase of string theory does also satisfy the swampland criteria.

Turning the logic around and using the swampland conjectures as a guide to the high temperature phase, we conclude that a purely topological theory without any propagating degrees of freedom might be too simple.

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