Klebanov-Strassler black hole

Alex Buchel

Department of Applied Mathematics,
Department of Physics and Astronomy
University of Western Ontario
London, Ontario N6A 5B7, Canada;
Perimeter Institute for Theoretical Physics
Waterloo, Ontario N2J 2W9, Canada

Abstract

We construct a black hole solution on warped deformed conifold in type IIB supergravity with fluxes. The black hole has translationary invariant horizon and is a holographic dual to a thermal homogeneous and isotropic state of a cascading $SU(K+P) \times SU(K)$ $\mathcal{N} = 1$ supersymmetric gauge theory with spontaneously broken chiral symmetry. We discuss thermal properties of the new black hole solutions. We comment on implications of the new black hole solutions for the landscape of KKLT de Sitter vacua in string theory.

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1 Introduction

A conifold is a complex 3-dimensional manifold described by the following equation in \( \mathbb{C}^4 \),

\[
\sum_{n=1}^{4} z_n^2 = 0.
\]

(1.1)

Owing to explicitly known Ricci-flat metric [1], the conifold appeared prominently in string theory and supergravity:

- In a holographic context [2], placing a large number \( K \) of D3 branes at the tip of the (singular) conifold realizes a duality for \( \mathcal{N} = 1 \) \( SU(K) \times SU(K) \) superconformal gauge theory, also known as Klebanov-Witten gauge theory [3]. Further wrapping \( P \) D5 branes on a 2-cycle of a conifold realizes a correspondence with \( \mathcal{N} = 1 \) \( SU(K + P) \times SU(K) \) Klebanov-Strassler (KS) gauge theory [4]. KS gauge theory undergoes an infinite sequence - a cascade - of self-similar Seiberg duality [5] transformations in the UV; it confines with the spontaneous chiral symmetry breaking in the IR.

- In [6] (GKP) it was pointed out that type IIB string theory compactified on warped throat geometries with fluxes (local version of which is precisely that of KS gauge theory gravitational dual) produces no-scale \( \mathcal{N} = 1 \) supersymmetric Minkowski vacua, naturally generating large hierarchies of physical scales. GKP compactifications fix all complex structure moduli, leaving at least a single Kähler modulus (an overall volume of a compact 6-dimensional manifold) unfixed. KKLT [7] further argued that non-perturbative corrections in GKP compactifications fix the overall volume Kähler
modulus leading to $\mathcal{N} = 1$ SUSY preserving $AdS_4$ vacua. Adding $D3$ to compactifications lifts AdS vacua to de Sitter. KKLT construction has been taken as the primary evidence for a landscape of de Sitter vacua in string theory.

In this paper we present new results regarding black hole solutions in type IIB supergravity on warped deformed conifold with fluxes. Our results are relevant both for holography and the landscape:

- Black hole geometries in holography represent gravitational dual to thermal states of the boundary gauge theories [8]. Although the chiral symmetry is spontaneously broken in the vacuum of KS gauge theory, it is expected to be restored at sufficiently high temperature. This predicts the existence of black hole solutions dual to thermal states of the cascading gauge theory plasma with unbroken chiral symmetry [9]. Such black holes resolve the singularity of the Klebanov-Tseytlin [10] geometry [9, 11, 12]. KS gauge theory confines in the infrared — the holographic dual of this (first-order) phase transition was established in [13]. ABK black hole solutions represent the gravitational dual to deconfined thermal states of KS gauge theory with unbroken chiral symmetry. These black holes cease to exist below some critical temperature $T < T_u$ [15], where they join a perturbatively unstable branch (with a negative specific heat and condensation of the hydrodynamic sound modes). ABK black holes are also perturbatively unstable for $T < T_{\chi_{SB}} = 1.00869(0) T_u$ towards development of the chiral symmetry breaking ($\chi_{SB}$) condensates in KS gauge theory plasma [16]. The end point of the $\chi_{SB}$ breaking instability in ABK black holes would produce Klebanov-Strassler black hole — gravitational backgrounds dual to homogeneous and isotropic thermal deconfined states of KS gauge theory plasma with spontaneously broken chiral symmetry. Until this work no such black hole was constructed. Thermo-dynamics of KS black holes will be discussed in section 2.

- The most controversial aspect of the KKLT construction of de Sitter vacua is the uplift of $AdS_4$ non-perturbative GKP vacua due to $D3$ branes. The backreacted solutions corresponding to smeared $D3$ branes at the tip of the KS solution were argued to be singular [17, 20]. Because the singularity is localized, it must

\[ ^1 \text{To understand the thermodynamics of chirally symmetric states of the cascading gauge theory plasma it was important to understand the holographic renormalization of the theory [14].} \]

\[ ^2 \text{There are incorrect claims in the literature regarding construction of black holes on warped deformed conifold with fluxes.} \]
be possible to study it in the local geometry of the GKP background — the noncompact gravitational dual to KS gauge theory. The following strategy was proposed in [21]: if the singularity due to $D^3$ branes at the conifold is physical, it should be possible to shield it with a horizon [22]: i.e., there must exist a black hole solution on the conifold that carries a negative $D^3$ brane charge at the horizon. Black holes with negative $D^3$ brane horizon charge have not been found for a conifold with an unbroken $U(1)$ symmetry (a gravitational dual to a chiral symmetry in KS gauge theory); neither were found black holes with negative charge where this $U(1)$ symmetry is broken explicitly [21]. KKLT construction requires warped deformed conifolds with spontaneous symmetry breaking $U(1) \supset \mathbb{Z}_2$ — thus, one needs to search for negative $D^3$ brane charge KS black holes. We report on this in section 3.

In the next two sections we present results of relevance to cascading gauge theory holography and to the landscape of KKLT de Sitter vacua, omitting all the technical details. Reader interested in technical details should consult [16].

2 Holography: phases of the cascading gauge theory

In our review of the cascading gauge theory and the thermodynamics of its chiral symmetric states we closely follow [16].

Klebanov-Strassler gauge theory is $\mathcal{N} = 1$ four-dimensional supersymmetric $SU(K+P) \times SU(K)$ gauge theory with two chiral superfields $A_1, A_2$ in the $(K+P, K)$ representation, and two fields $B_1, B_2$ in the $(K+P, K)$. This gauge theory has two gauge couplings $g_1, g_2$ associated with two gauge group factors, and a quartic superpotential

$$W \sim \text{Tr} (A_i B_j A_k B_\ell) \epsilon^{ik} \epsilon^{j\ell}. \quad (2.1)$$

When $P = 0$ above theory flows in the infrared to a superconformal fixed point, commonly referred to as Klebanov-Witten theory. At the IR fixed point KW gauge theory is strongly coupled — the superconformal symmetry together with $SU(2) \times SU(2) \times U(1)$ global symmetry of the theory implies that anomalous dimensions of chiral superfields $\gamma(A_i) = \gamma(B_i) = -\frac{1}{4}$, i.e., non-perturbatively large.

When $P \neq 0$, conformal invariance of the above $SU(K+P) \times SU(K)$ gauge theory is broken. It is useful to consider an effective description of this theory at energy scale
\( \mu \) with perturbative couplings \( g_i(\mu) \ll 1 \). It is straightforward to evaluate NSVZ beta-functions for the gauge couplings. One finds that while the sum of the gauge couplings does not run
\[
\frac{d}{d \ln \mu} \left( \frac{\pi}{g_s} = \frac{4\pi}{g_1^2(\mu)} + \frac{4\pi}{g_2^2(\mu)} \right) = 0, \tag{2.2}
\]
the difference between the two couplings is
\[
\frac{4\pi}{g_2^2(\mu)} - \frac{4\pi}{g_1^2(\mu)} \sim P \left[ 3 + 2(1 - \gamma_{ij}) \right] \ln \frac{\mu}{\Lambda}, \tag{2.3}
\]
where \( \Lambda \) is the strong coupling scale of the theory and \( \gamma_{ij} \) is an anomalous dimension of operators \( \text{Tr} A_i B_j \). Given (2.3) and (2.2) it is clear that the effective weakly coupled description of \( SU(K + P) \times SU(K) \) gauge theory can be valid only in a finite-width energy band centered about \( \mu \) scale: extending effective description both to the UV and to the IR one necessarily encounters strong coupling in one or the other gauge group factor. To extend the theory past the strongly coupled region(s) one must perform a Seiberg duality. In KS gauge theory a Seiberg duality transformation is a self-similarity transformation of the effective description so that \( K \to K - P \) as one flows to the IR, or \( K \to K + P \) as the energy increases. Thus, extension of the effective \( SU(K + P) \times SU(K) \) description to all energy scales involves a cascade of Seiberg dualities where the rank of the gauge group changes with energy according to
\[
K = K(\mu) \sim 2P^2 \ln \frac{\mu}{\Lambda}, \tag{2.4}
\]
at least as \( \mu \gg \Lambda \). Although there are infinitely many duality cascade steps in the UV, there is only a finite number of duality transformations as one flows to the IR (from a given scale \( \mu \)). The space of vacua of a generic cascading gauge theory was studied in details in [23]. When \( K(\mu) \) is an integer multiple of \( P \), the cascading gauge theory confines in the infrared with a spontaneous breaking of the chiral symmetry.

The thermal phase digram of homogeneous and isotropic states in cascading gauge theory plasma represents competition between three phases:

- (A): confined phase with spontaneously broken chiral symmetry;
- (B): deconfined chirally symmetric phase;
- (C): deconfined phase with spontaneously broken chiral symmetry.

Correspondingly, in a dual gravitational description we have:
• (Ah): thermal KS geometry, i.e., Klebanov-Strassler vacuum solution with periodically identified Euclidean time direction,
\[ t_E \sim t_E + \frac{1}{T}. \] (2.5)

In this phase all the thermodynamic potentials vanish: the free energy density \( \mathcal{F} \), the energy density \( \mathcal{E} \) and the entropy density \( s \). There are nonvanishing condensates of two dimension-3 operators (dual to chiral symmetry breaking gaugino condensates of both gauge group factors), and a condensate of a dimension-6 operator \[16\].

• (Bh): Klebanov-Tseytlin black hole. In this phase we have nonvanishing \( \{\mathcal{F}, \mathcal{E}, s\} \). There are nonvanishing condensates of two dimension-4 operators, a dimension-6 operator and a dimension-8 operator \[13\]. Condensates of the chiral symmetry breaking operators vanish.

• (Ch): Klebanov-Strassler black hole. In this phase we have nonvanishing \( \{\mathcal{F}, \mathcal{E}, s\} \). In addition to the condensates present in (Bh), we have condensates of a pair of chiral symmetry breaking dimension-3 operators (as in (Ah)) and an additional condensate of a dimension-7 operator (also breaking the chiral symmetry) \[16\].

Phase transition between \( A \leftrightarrow B \) is of the first-order \[13\].

We now turn to a detailed discussion of the phase diagram of the cascading gauge theory plasma. At temperatures \( T \gg \Lambda \) the cascading plasma is in the deconfined phase with an unbroken chiral symmetry (B) \[9\] \[11\] \[12\]. Here, the temperature-dependent effective rank \( K(T) \) of the cascading theory is large, compare to \( P \) \[14\]:
\[ \frac{K(T)}{P^2} = \frac{1}{2} \ln \left( \frac{64\pi^4}{81} \times \frac{sT}{\Lambda^4} \right) \implies \frac{K(T)}{P^2} \approx 2 \ln \frac{T}{\Lambda}, \quad T \gg \Lambda. \] (2.6)

To leading order at higher temperature, the pressure \( P = -\mathcal{F} \) and the energy density \( \mathcal{E} \) are given by \[14\]
\[ \frac{\mathcal{P}}{sT} = \frac{1}{4} \left( 1 - \frac{P^2}{K(T)^2} + \mathcal{O} \left( \frac{P^4}{K(T)^2} \right) \right), \]
\[ \frac{\mathcal{E}}{sT} = \frac{3}{4} \left( 1 + \frac{1}{3} \frac{P^2}{K(T)^2} + \mathcal{O} \left( \frac{P^4}{K(T)^2} \right) \right). \] (2.7)

At low temperature/energy density we need to distinguish canonical (see fig. 1) and microcanonical (see fig. 2) ensembles.
Figure 1: Phase diagram of the cascading gauge theory plasma in canonical ensemble. Solid blue line represents deconfined, chirally symmetric phase of the theory — the gravitational dual to these thermal states is a Klebanov-Tseytlin black hole. Solid red lines represent deconfined phase of the theory with the spontaneously broken chiral symmetry — the gravitational dual to these thermal states is a Klebanov-Strassler black hole. Left panel shows the free energy density versus the temperature of the symmetric and the symmetry broken phases. The vertical dashed green line indicates the first-order confinement-deconfinement phase transition; the vertical dashed orange line indicates the onset of the perturbative instability of the chirally symmetric phase; the vertical dashed black line indicates the onset of the instability of sound waves in chirally symmetric phase of the cascading plasma. Right panel shows a critical behavior of one of the dimension-3 condensates in the chiral symmetry broken phase of the cascading plasma.

2.1 Canonical ensemble

The deconfined chirally symmetry phase (B) (represented by a solid blue curve) extends to temperature \[ T_c = 0.6141111(3)\Lambda, \] below which its free energy density \[ \tilde{F} = \frac{128\pi^4}{81P^4} F \] becomes positive. This signifies the first-order phase transition to the confined phase with the spontaneously broken chiral symmetry (A). This phase transition is denoted by a vertical dashed green line. Since \( A \leftrightarrow B \) phase transition proceeds via bubble
nucleation, it is non-perturbative. At temperature
\[ T_{\chi_{SB}} = 0.54195(5) \Lambda, \]
the meta-stable phase (B) becomes perturbatively unstable due to chiral symmetry breaking fluctuations\(^{[16]}\). This instability is denoted by a vertical dashed orange line. At temperature
\[ T_u = 0.537286 \Lambda, \]
the phase (B) terminates joining a perturbatively unstable branch of the theory with a negative specific heat\(^{[15]}\). The branch with a negative specific heat has dynamical instability leading to a breakdown of spatial homogeneity in plasma: the sound waves are unstable\(^{[24]}\). The terminal temperature \( T_u \) is denoted by a vertical dashed black line. Note the hierarchy of critical temperatures of homogeneous and isotropic thermal states in cascading plasma:
\[ T_u < T_{\chi_{SB}} < T_c. \]

A natural expectation is that the deconfined phase with spontaneously broken chiral symmetry — the phase (C) — should bifurcate from (B) at \( T = T_{\chi_{SB}} \) where the fluctuations associated with this symmetry breaking become unstable. The end point of the instability for \( T < T_{\chi_{SB}} \) would be Klebanov-Strassler black hole. Until this work, the searches for the KSBH were unsuccessful. The right panel of fig. 1 provides a reason\(^{[3]}\) although the \( \chi_{SB} \) fluctuations are unstable for \( T \leq T_{\chi_{SB}} \), the KSBH exists only for \( T \geq T_{\chi_{SB}} \). In other words, the KSBH is in a class of exotic black holes originally identified in\(^{[25]}\). The phase (C) of the cascading gauge theory plasma is denoted by a solid red curve: it has a higher free energy density than the phase (B) at the corresponding temperature and thus never dominates in macrocanonical ensemble.

### 2.2 Microcanonical ensemble

Microcanonical ensemble is relevant for dynamical questions (thermalization and equilibration) of gauge theory plasma. Fig. 2 presents the phase diagram of the cascading gauge theory in microcanonical ensemble. The solid blue curve denotes phase (B), and the solid red curve denotes phase (C). Similar to\(^{2}(2.9)\) we introduced
\[ \hat{E} = \frac{128\pi^4}{81P^4} E, \quad \hat{s} = \frac{128\pi^4}{81P^4} s. \]

\(^3\)Over the years, the author was searching for the KSBH at \( T < T_{\chi_{SB}} \).

\(^4\)See also\(^{[26]}\).
Figure 2: Phase diagram of the cascading gauge theory plasma in microcanonical ensemble. Solid blue lines represent deconfined, chirally symmetric phase of the theory — the gravitational dual to these thermal states is a Klebanov-Tseytlin black hole. Solid red lines represent deconfined phase of the theory with the spontaneously broken chiral symmetry — the gravitational dual to these thermal states is a Klebanov-Strassler black hole. Left panel shows the entropy density versus the energy density of the symmetric and the symmetry broken phases. The vertical dashed orange line indicates the onset of the perturbative instability of the chirally symmetric phase; the vertical dashed black line indicates the onset of the instability of the sound waves in chirally symmetric phase of the cascading plasma. Right panel shows the speed of sound waves squared in symmetric and symmetry broken phases of the cascading plasma.

A vertical orange line
\[ \hat{\mathcal{E}}_{\chi \text{SB}} = 1.270093(1) \Lambda^4 \]  
indicates the onset of the chiral symmetry breaking instability [16]. Notice that here the phase (C) exists for \( \mathcal{E} \leq \mathcal{E}_{\chi \text{SB}} \) and dominates over the phase (B). In other words, insisting on homogeneous and isotropic evolution, the KSBH is the end point of the perturbative instability of the KTBH at sufficiently low energy densities. Such a phenomenon in a context of exotic black holes was identified in [27].

The KSBH is both thermodynamically and dynamically unstable. The right panel of fig. 2 presents the speed of sound waves as a function of the energy density for (B) (blue curve) and (C) (red curve) phases of the cascading gauge theory plasma. A vertical black line
\[ \hat{\mathcal{E}}_u = 0.723488 \Lambda^4 \]  
indicates the onset of the perturbative instability in phase (B) associates with the breaking of the translational invariance due to the condensation of the hydrodynamic
sound waves \[15\].

3 Landscape: KKLT de Sitter vacua in string theory

Following \[21\], we compute the Maxwell D3-brane charge of the conifold black hole horizons $Q^D_3$. A negative value of the horizon charge would indicate that the anti-D3 brane singularity is physical, according to classification \[22\]. The results of the computations are presented in fig. 3. The solid blue curve represents the Maxwell charge of KT BH, and the solid red curve represents the charge of the KSBH. Both charges are never negative; for $\mathcal{E} < \mathcal{E}_{\chi SB}$, when the KSBH has a higher entropy then the corresponding energy density KT BH,

\[
Q^D_3\bigg|_{KSBH} > Q^D_3\bigg|_{KTBH}.
\]  

(3.1)
4 Conclusion

In this paper we reported on the physical properties of Klebanov-Strassler black holes on warped deformed conifold with fluxes. These black holes are important for understanding of the holographic correspondence for the confining KS gauge theory [4]. They were also long-sought objects in the context of KKLT de Sitter vacua constructions in string theory [7].

We established the existence of KSBHs. We determined that these black holes have a negative specific heat, and are dynamically unstable due to the hydrodynamic sound wave fluctuations (breaking the homogeneity of the horizon).

KSBHs realize the homogeneous and isotropic deconfined phase of the cascading gauge theory plasma with spontaneously broken chiral symmetry. While the corresponding phase never dominates in canonical ensemble, it has a higher entropy density compare to the homogeneous and isotropic phase with unbroken chiral symmetry at the corresponding energy density (the gravitational dual to KTBHs) below some critical energy density, \( \mathcal{E} < \mathcal{E}_{\chi SB} \). Of course, since ultimately the homogeneity assumption is not valid (KSBHs are dynamically unstable), the equilibrium states of the cascading gauge theory below \( \mathcal{E}_{\chi SB} \) remain unknown — for sure they can not be homogeneous and isotropic.

Finally, we demonstrated that KSBHs can not shield (conjectured) anti-D3 brane singularity on the warped deformed conifold with fluxes.

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