



# Black holes in $\omega$ -deformed gauged $N = 8$ supergravity



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## ABSTRACT

Motivated by the recently found 4-dimensional  $\omega$ -deformed gauged supergravity, we investigate the black hole solutions within the single scalar field consistent truncations of this theory. We construct black hole solutions that have spherical, toroidal, and hyperbolic horizon topologies. The scalar field is regular everywhere outside the curvature singularity and the stress–energy tensor satisfies the null energy condition. When the parameter  $\omega$  does not vanish, there is a degeneracy in the spectrum of black hole solutions for boundary conditions that preserve the asymptotic Anti-de Sitter symmetries. These boundary conditions correspond to multi-trace deformations in the dual field theory.

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## 1. Introduction

Supergravity is an effective theory for string theory at low energies compared to the string energy scale. The Anti-de Sitter/Conformal Field Theory (AdS/CFT) duality [1] provides a remarkable tool for extracting information about strongly coupled gauge theories (in  $d$  dimensions) from a dual supergravity description (in  $d + 1$  dimensions). Within the AdS/CFT duality, the radial coordinate plays the role of the energy scale and so the bulk geometry has a nice interpretation as a renormalization group (RG) flow of the dual field theory [2] (see, also, the review [3]).

The gauged  $N = 8$  supergravity in four dimensions is the maximal gauged supergravity with spins lower or equal than two and it can be obtained by Kaluza–Klein reduction of 11-dimensional supergravity [4] on a 7-dimensional sphere [5]. In the holographic context, much of the interest on the 4-dimensional gauged supergravities comes from the utility of the ABJM model [6] in testing various strongly coupled phenomena in condensed matter physics.

What came as a surprise, recently, is the existence of a continuous one-parameter family of inequivalent maximally supersymmetric gauged supergravities [7] (the critical points were extensively studied in [8]). These theories are characterized by one ‘angular parameter’,  $\omega$ , that generates an electric–magnetic duality transformation prior to selecting the  $SO(8)$  gauging, and they are referred to as  $\omega$ -deformed gauged  $N = 8$  supergravity. At the practical level, the  $\omega$  parameter can be introduced in the maximal

gauged supergravity by multiplying the 56-bein,  $\mathcal{V}$ , by a diagonal matrix containing either  $e^{i\omega}$  or  $e^{-i\omega}$  [9]. As the moduli potential is a non-linear function in terms of  $\mathcal{V}$ , the parameter  $\omega$  appears explicitly in the potential. It is worth noting that, long time ago, this kind of ‘angles’ was introduced in  $N = 4$  gauged supergravities in [10].

In this Letter, we construct *exact* domain wall (see, also, [11, 12]) and black hole solutions in one scalar field consistent truncations of  $\omega$ -deformed theories [11], which correspond to RG flows of 3-dimensional dual quantum field theories at zero and finite temperatures, respectively. We show that the null energy condition (which is the relevant energy condition in AdS) is satisfied and construct the  $c$ -function using the gravity side of the duality.

For  $\omega = 0$ , we obtain domain wall solutions and explicitly write down the corresponding superpotential. When  $\omega$  does not vanish, there is a degeneracy in the spectrum of black hole solutions. Since the black holes are interpreted as thermal states in the dual theory, this degeneracy may seem puzzling at first sight. However, this can be understood because the solutions correspond to different boundary conditions which in turn correspond to different deformations of the dual field theory. The degeneracy in the spectrum of solutions can then be understood as a sign degeneracy in the AdS invariant boundary conditions due to a change in the sign of the scalar field. It is also interesting to note that the no-hair conjecture in asymptotically AdS spacetimes of [14] have as a hypothesis the reality of the superpotential. So, it does not apply to the  $\omega$ -deformed theory, as the superpotential is generically complex [15].

Since the coordinates in which the black hole solutions are constructed are not very intuitive, we will explicitly provide the

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change of coordinates that asymptotically puts the metric in the canonical form [16]. In this system of coordinates, it is easy to observe that the scalar field satisfies mixed boundary conditions that correspond to multi-trace deformations in the dual theory [17].

### 2. Moduli potential

The single scalar field truncations of the  $\omega$ -deformed theory were found in [11]. The action is the one of Einstein gravity minimally coupled to a scalar field

$$I[g_{\mu\nu}, \varphi] = \frac{1}{16\pi G} \int d^4x \sqrt{-g} \left[ R - \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi - V_n(\varphi) \right] \quad (1)$$

where  $G$  is the Newton constant. The moduli potentials are parameterized by an integer,  $n = \{1..7\}$ , and they contain a non-trivial deformation parameter,  $\omega$ , when  $n$  is odd [11].<sup>1</sup>

We found that all the truncations are contained in a more general moduli potential

$$W_\nu(\varphi) = \frac{2\alpha}{\nu^2} \left[ \frac{\nu-1}{\nu+2} \sinh(\varphi l_\nu(\nu+1)) - \frac{\nu+1}{\nu-2} \sinh(\varphi l_\nu(\nu-1)) + 4 \frac{\nu^2-1}{\nu^2-4} \sinh(\varphi l_\nu) \right] - \frac{(\nu^2-4)}{l^2 \nu^2} \left[ \frac{\nu-1}{\nu+2} \exp(-\varphi l_\nu(\nu+1)) + \frac{\nu+1}{\nu-2} \exp(\varphi l_\nu(\nu-1)) + 4 \frac{\nu^2-1}{\nu^2-4} \exp(-\varphi l_\nu) \right] \quad (2)$$

where  $l$  is the AdS radius and  $l_\nu^{-2} = (\nu^2 - 1)$ . To obtain the potentials  $V_n(\varphi)$  from the general potential  $W_\nu(\varphi)$ , we have to relate the parameters  $\omega$  and  $g$  to  $\alpha$  and the AdS radius,  $l$ , in the following way:  $\sin(\omega)^2 = \frac{\alpha S_n}{g^2}$ ,  $g = \frac{1}{\sqrt{2}l}$ , where for odd  $n$  we define

$S_1 = \frac{9}{40}$ ,  $S_3 = -\frac{1}{24}$ ,  $S_5 = \frac{1}{24} + \frac{g^2}{\alpha}$ , and  $S_7 = -\frac{9}{40} + \frac{g^2}{\alpha}$ . Then, by using the relations between parameters, it is straightforward to show that

$$\begin{aligned} V_1(\varphi) &= W_{\frac{4}{3}}(\varphi) = V_7(-\varphi), \\ V_2(\varphi) &= W_2(\varphi)|_{\alpha=0} = V_6(\varphi) \\ V_3(\varphi) &= W_4(\varphi) = V_5(-\varphi), \quad V_4(\varphi) = W_\infty(\varphi)|_{\alpha=0} \end{aligned} \quad (3)$$

The scalar field potential (2) has a negative maximum and so AdS is a solution.<sup>2</sup> More precisely, the AdS vacuum can be obtained for  $\varphi = 0$

$$W_\nu(\varphi) = -\frac{3}{l^2}, \quad \frac{dW_\nu(\varphi)}{d\varphi} = 0, \quad \frac{d^2W_\nu(\varphi)}{d\varphi^2} = -\frac{2}{l^2} \quad (5)$$

Indeed, the small scalar fluctuations are tachyonic,  $m^2 < 0$ , but the mass is still above the Breitenlohner–Freedman (BF) bound,  $m_{BF}^2 = -\frac{9}{4l^2}$  and so this vacuum is perturbatively stable [18]. Moreover, since the mass is in the range  $m_{BF}^2 < m^2 < m_{BF}^2 + l^{-2}$ , both modes of the scalar field are normalizable [19]. Therefore, these theories are suitable in the context of AdS/CFT duality with boundary conditions corresponding to the presence of multi-trace operators in the dual field theory [17].

We would like to point out that it can be shown that for these one scalar field consistent truncations, the range of the parameter

<sup>1</sup>  $V_1(\varphi)$  was also discussed and independently derived in [9].  
<sup>2</sup> Obviously, this is the critical point that preserves  $N = 8$  supersymmetry and the group  $SO(8)$ , however there exist also additional non-supersymmetric critical points [8].

is  $\omega \in [0, \frac{\pi}{4}]$  [11] rather than  $\omega \in [0, \frac{\pi}{8}]$  like in the case of the full theory [7]. The reason behind this fact is that the fields for which the identifications should be done in order to reduce the range of  $\omega$  have been truncated out. When  $\alpha = 0$  these single scalar field truncations correspond to different cases of the four dimensional potentials of [13].

### 3. Exact black hole solutions

In this section, we investigate the existence of event horizons in terms of the parameters of the potential. Similar techniques to find hairy black holes were used in [20–22] (see, also, [23,24]) where the details of the construction can be found. In what follows, we present the solutions, discuss some of their properties, and explicitly check the existence of the horizon.

The black holes in AdS can have non-spherical horizon topology and the solutions relevant to  $\omega$ -deformed gauge supergravity can be written in a compact form as

$$ds^2 = \Omega(x) \left[ -f(x) dt^2 + \frac{\eta^2 dx^2}{f(x)} + d\Sigma_k \right] \quad (6)$$

$$\varphi = l_\nu^{-1} \ln(x), \quad \Omega(x) = \frac{\nu^2 x^{\nu-1}}{\eta^2 (x^\nu - 1)^2} \quad (7)$$

$$f(x) = \frac{x^{2-\nu} (x^\nu - 1)^2 \eta^2 k}{\nu^2} + \left[ \frac{1}{\nu^2 - 4} - \frac{x^2}{\nu^2} \left( 1 + \frac{x^{-\nu}}{\nu - 2} - \frac{x^\nu}{\nu + 2} \right) \right] \alpha + \frac{1}{l^2} \quad (8)$$

Here,  $d\Sigma_k$  is a two-manifold of constant curvature  $k = \pm 1$  or 0 and the corresponding black holes can have spherical, toroidal or hyperbolic horizon topologies. The parameter  $\eta$  is in fact the integration constant and the reason it appears in a non-standard form in the metric is because the computations simplify when a dimensionless radial coordinate,  $x$ , is used.

The configuration and scalar potential are invariant under the change  $\nu \rightarrow -\nu$ , therefore from now on we consider  $\nu \geq 1$ . We also note that conformal infinity is at  $x = 1$ . There are two branches of solutions depending on whether  $x \in (0, 1)$  or  $x \in (1, \infty)$ . The metric is regular for any value of  $x \neq 0$  and  $x \neq \infty$  as can be seen from the introduction of advanced and retarded coordinates  $u_\pm = t \mp \int \frac{\eta}{f(x)} dx$ . The scalar field and metric are singular at  $x = 0$  and  $x = \infty$  but, as we will see, these singularities can be covered by event horizons.

Interestingly, it is straightforward to observe in our parameterization that, with this scalar field potential, the theory also admits asymptotically flat hairy black holes [22,25,26] (when  $l^{-2} = 0$ ) that could, in principle, be related to asymptotically flat gauged supergravity [27]. Solutions of the form (6)–(8) in four dimensions were originally found in [20] and discussed in different contexts in the literature, [23,28–30]. Indeed, these solutions exist for any value of  $\nu$  and  $\alpha$  although we shall focus below only on the values related to gauged supergravity.

For even  $n$  and vanishing  $\alpha$  the situation is well known and has been thoroughly studied. When  $n$  is even, there exist exact neutral asymptotically locally AdS black holes only for  $k = -1$  [31]. In what follows, we are going to carefully analyze the existence of event horizons for  $k = 1$  and comment also on the case  $k = 0$ . Since the solutions are static, the horizon is localized at the place where the  $g_{tt}$  component of the metric vanishes. That is, there is an  $x_+$  such that  $f(x_+) = 0$ . As expected, since the solution is asymptotically AdS, we obtain at the boundary, where the scalar field vanishes,  $f(x = 1) = l^{-2}$ . Furthermore, it is straightforward to obtain the (non-trivial) solution of the equation  $\frac{df}{dx} = 0$ , and by

analyzing the second derivative of  $f$  we obtain that this point is a maximum. Therefore, the function  $f$  has at most one zero in any physically relevant interval. Since  $f$  is positive at the boundary, the existence of a black hole horizon is ensured if  $f$  is negative near the singularities. Therefore, the necessary and sufficient condition for the existence of a horizon are

$$\begin{aligned} \nu < 2, x < 1 &\implies \frac{\alpha}{\nu^2 - 4} + \frac{1}{l^2} < 0 \\ \nu < 2, x > 1 &\implies \left( \eta^2 + \frac{\alpha}{\nu + 2} \right) < 0 \end{aligned} \quad (9)$$

$$\begin{aligned} \nu > 2, x < 1 &\implies \left( \eta^2 - \frac{\alpha}{\nu - 2} \right) < 0 \\ \nu > 2, x > 1 &\implies \left( \eta^2 + \frac{\alpha}{\nu + 2} \right) < 0 \end{aligned} \quad (10)$$

Therefore, there should exist hairy black holes in an open set of the parameter space when at least one of the relations above are satisfied.

Finally, we want to point out that the symmetries of these theories (3) and (4) allow to map the  $n = 3$  into the  $n = 5$  theory by making the field redefinition  $\varphi \rightarrow -\varphi$  and letting the geometry invariant.

### 3.1. Spherically symmetric black holes

We are going to consider the cases of non-trivial  $\alpha$ , namely when  $n$  is odd.

- When  $n = 1$  then  $\alpha = \frac{\sin(\omega)^2}{l^2} \frac{20}{9}$  and  $\nu = \frac{4}{3}$ . Then, the condition (9) shows that there are no spherically symmetric black holes.
- When  $n = 3$  then  $\alpha = -\frac{12 \sin(\omega)^2}{l^2}$  and  $\nu = 4$ . Then, the condition (10) shows that there are black holes when  $x > 1$  and  $\sin(\omega)^2 > \frac{l^2 \eta^2}{2}$ .
- When  $n = 5$  then  $\sin(\omega)^2 = \frac{\alpha l^2}{12} + 1 \implies \cos(\omega)^2 = -\frac{\alpha l^2}{12}$  and  $\nu = 4$ . Then, the condition (10) shows that there are black holes when  $x > 1$  and  $\cos(\omega)^2 > \frac{l^2 \eta^2}{2}$ .
- When  $n = 7$  then  $\sin(\omega)^2 = -\frac{9\alpha l^2}{20} + 1 \implies \cos(\omega)^2 = \frac{9\alpha l^2}{20}$  and  $\nu = \frac{4}{3}$ . Then, the condition (9) shows that there are no spherically symmetric black holes.

### 3.2. Planar black holes

For  $k = 0$ , the condition that  $f(x)$  becomes negative near the singularity is

$$\begin{aligned} \nu < 2, x < 1 &\implies \frac{\alpha}{\nu^2 - 4} + \frac{1}{l^2} < 0 \\ \nu < 2, x > 1 &\implies \frac{\alpha}{\nu + 2} < 0 \end{aligned} \quad (11)$$

$$\begin{aligned} \nu > 2, x < 1 &\implies -\frac{\alpha}{\nu - 2} < 0 \\ \nu > 2, x > 1 &\implies \frac{\alpha}{\nu + 2} < 0 \end{aligned} \quad (12)$$

As in the previous case, there are no black holes when  $n = 1$  or 7.

- When  $n = 3$  then condition (10) shows that there are black holes when  $x > 1$  and for  $\forall \omega \neq 0$ .
- When  $n = 5$  then condition (10) shows that there are black holes when  $x > 1$  and for  $\forall \omega$ .

When  $k = 0$  and  $\alpha = 0$ , we also obtain domain wall solutions:

$$ds^2 = l^2 \Omega(x) [-dt^2 + \eta^2 dx^2 + dy^2 + dz^2] \quad (13)$$

$$\varphi = l_\nu^{-1} \ln(x), \quad \Omega(x) = \frac{\nu^2 x^{\nu-1}}{\eta^2 (x^\nu - 1)^2} \quad (14)$$

In this case the potential can be explicitly written in terms of a superpotential

$$W_\nu(\varphi) = 2 \left( \frac{dP(\varphi)}{d\varphi} \right)^2 - \frac{3}{2} P(\varphi)^2 \quad (15)$$

where

$$P(\varphi) = l^{-1} \left[ \frac{(\nu + 1)}{\nu} e^{(\frac{\nu-1}{2})\varphi l_\nu} + \frac{(\nu - 1)}{\nu} e^{-(\frac{\nu+1}{2})\varphi l_\nu} \right] \quad (16)$$

## 4. Boundary conditions

One important question we would like to address in this section is if the boundary conditions for our solutions are compatible with the AdS/CFT duality. The mixed boundary conditions compatible with AdS/CFT duality were interpreted as a multi-trace deformation of the boundary CFT [17].

Let us discuss the branch  $x \in (1, \infty)$  for which the scalar field is positively defined. We change the coordinates so that the function in front of the transversal section,  $d\Sigma_k$ , has the following fall-off:

$$\Omega(x) = r^2 + O(r^{-3}) \quad (17)$$

By using the change of coordinates<sup>3</sup>

$$x = 1 + \frac{1}{\eta r} - \frac{(\nu^2 - 1)}{24\eta^3 r^3} \left[ 1 - \frac{1}{\eta r} - \frac{9(\nu^2 - 9)}{80\eta^2 r^2} \right] \quad (18)$$

we obtain the following asymptotic expansions for the metric functions:

$$-g_{tt} = f(x)\Omega(x) = \frac{r^2}{l^2} + k + \frac{\alpha + 3\eta^2 k}{3\eta^3 r} + O(r^{-3}) \quad (19)$$

$$\begin{aligned} g_{rr} = \frac{\Omega(x)\eta^2}{f(x)} \left( \frac{dx}{dr} \right)^2 &= \frac{l^2}{r^2} - \frac{l^4}{r^4} - \frac{l^2(\nu^2 - 1)}{4\eta^2 r^4} \\ &+ \frac{l^2(3kl^2\eta^2 + l^2\alpha - \nu^2 + 1)}{3\eta^3 r^5} + O(r^{-6}) \end{aligned} \quad (20)$$

Therefore, the asymptotic behavior of our exact black hole solutions fits very well in the general analysis of [16] and, consequently, we expect that the boundary conditions for the scalar correspond to a multi-trace deformation in the dual field theory [17]. The asymptotic expansion of the scalar field is

$$l_\nu \varphi = \frac{1}{\eta r} - \frac{1}{2\eta^2 r^2} - \frac{\nu^2 - 9}{24\eta^3 r^3} + O(r^{-4}) \quad (21)$$

and by comparing with the scalar field in the standard notation of [16]

$$\varphi = \frac{a}{r^{\lambda_-}} + \frac{b}{r^{\lambda_+}} + \dots \quad \lambda_- < \lambda_+ \quad (22)$$

we obtain  $\lambda_+ = 2\lambda_- = 2$ ,  $a = \frac{1}{\eta l_\nu}$ , and  $b = -\frac{1}{2\eta^2 l_\nu}$ . The scalar field introduces an asymptotic deviation at the order  $O(r^{-4})$  in (20) that is not present when the scalar field vanishes ( $\nu^2 = 1$ ).

Since  $b = -\frac{1}{2}a^2$ , these boundary conditions are, indeed, compatible with the canonical realization of the conformal symmetry at the boundary. However, when the scalar field is negative, the

<sup>3</sup> We assume that  $\eta > 0$  so that  $r \rightarrow \infty$  implies  $x \rightarrow 1$  from the right.

boundary conditions change and are of the form  $b = \frac{1}{2}a^2$  and there is a degeneracy in the spectrum of black hole solutions. These boundary conditions are invariant under the asymptotic AdS symmetries. Black holes are interpreted as thermal states in the dual field theory, and from this point of view,  $b$  is the expectation value of an operator of dimension two. The degeneracy is absent when  $\omega = 0$  because the black holes with  $n = 3$  and  $\varphi > 0$  do not exist in this case. It is also worth to remark that AdS invariant multi-trace deformation are not incompatible with supersymmetry, as was checked in some concrete examples in [32].

## 5. Discussion

In this Letter, we have investigated the existence of *exact* domain wall and black hole solutions in one scalar field consistent truncations of  $\omega$ -deformed  $N = 8$  SUGRA theories. The black holes only exist in the  $n = 3$  and  $n = 5$  theory, which are actually identical up to the field redefinition  $\varphi \rightarrow -\varphi$ . Let us therefore focus on the theory with  $n = 5$ . This theory has a hairy black hole with AdS invariant boundary conditions even when  $\omega = 0$  (because  $\alpha$  does not vanish in this case). The value of the scalar field is everywhere positive as it has no nodes. By turning on the deformation parameter, the theory has a new black hole solution that corresponds to a negative scalar field,  $\varphi < 0$ . From our discussion in Section 4, it is clear that this is equivalent to the fact that there are black holes for boundary conditions of the form  $b = \pm \frac{1}{2}a^2$ . Therefore, when  $\omega \neq 0$  there are *two* hairy black holes for each value of the parameters. Each hairy black hole is associated to a sign flip in the expectation value of a dimension 2 operator in the CFT. We would also like to remark that, to the best of our knowledge, the exact form of the black holes in undeformed gauged supergravity was unknown until now.

The existence of these black holes also implies that there is a richer phase diagram in the dual theory. Hairy black holes in  $AdS_4$  are physically relevant objects as they are an important ingredient of the AdS/Condensed Matter correspondence (see, e.g., [33] and the reference therein). As was originally pointed out in [34], there could also exist second order phase transitions between the hairy black holes and the hairless ones.

For completeness of the analysis, let us now show that the null energy condition is satisfied. Indeed, whenever there is a single minimally coupled scalar field to a diagonal metric it is possible to show that the null energy condition is satisfied. In an orthonormal frame  $e_{\mu}^a$ , the energy momentum tensor has the form  $T_{\mu\nu}e_a^{\mu}e_b^{\nu} = \text{diag}(\rho, p_1, p_2, p_2)$  where

$$\rho + p_1 = g_{xx}^{-1} \left( \frac{d\varphi}{dx} \right)^2 \geq 0, \quad \rho + p_2 = 0. \quad (23)$$

Since the null energy condition is valid for our solutions, there are no violations that lead to superluminal propagation and instabilities in the bulk [35] and so we expect that the boundary theory is well defined (see, e.g., [36]). Also, the c-function [37] can be also constructed<sup>4</sup>:

$$C(x) = C_0 \left( \frac{2\eta\Omega^{3/2}}{\Omega'} \right)^2 = C_0 \left[ \frac{x^{\frac{\nu+1}{2}}}{x^{\nu}(\nu+1) + (\nu-1)} \right]^2 \quad (24)$$

As expected, in the hairless limit when the neutral spherical black hole is obtained ( $\nu = 1$ ), the flow is trivial. We also notice that the c-function is completely determined by the conformal factor that is in agreement with the interpretation of the black hole as

a thermal state in the boundary theory. The c-function should remain unchanged when a finite temperature vacuum is excited in the same theory and so there should not be any dependence of the metric function  $f(x)$ .

In conclusion, the black hole solutions we have found correspond to mixed boundary conditions in AdS and can be interpreted within the AdS/CFT duality as RG flows of the dual field theory at finite temperature.

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<sup>4</sup> Similar examples are presented in [23,38] and we do not present the details here.

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