Supersymmetric Wormholes in String Theory

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We construct a large family of Euclidean supersymmetric wormhole solutions of type IIB supergravity which are asymptotically $\mathrm{AdS}_5 \times S^5$. The solutions are constructed using consistent truncation to maximally gauged supergravity in five dimensions which is further truncated to a four scalar model. Within this model we perform a full analytic classification of supersymmetric domain wall solutions with flat Euclidean domain wall slices. On each side of the wormhole, the solution asymptotes to AdS_5 dual to $\mathcal{N} = 4$ supersymmetric Yang-Mills deformed by a supersymmetric mass term.

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Introduction.—The Euclidean path integral for quantum gravity is an important topic of research and for low-dimensional theories such as JT gravity, has recently lead to many fruitful results; see, for instance, Ref. [1]. It has become clear that in low-dimensional theories it makes sense to sum over saddle points with different topologies [2]. Still, in higher dimensions for standard Einstein-Hilbert gravity (coupled to matter) the rules remains somewhat unclear. The story in higher dimensional gravity theories can be different from low-dimensional ones without leading to obvious inconsistencies.

In this regard, the role of wormholes, as possible saddle points of the path integral is still an important open problem [3]. The processes that involve wormholes pose puzzles for unitarity of the quantum system and nonfactorization of correlation functions in the holographic dual [4]. Moreover, the existence of wormholes indicates that probability amplitudes to produce or absorb baby universes are nontrivial which may lead to issues for the Swampland program [5,6].

To improve our understanding it is necessary to provide the embedding of higher dimensional Euclidean wormholes in string theory and AdS/CFT. In this way, various ideas regarding the semiclassical formulation of gravity can be put to a test. Research in this direction was initiated in some earlier works [7,8].

In order to construct Euclidean wormhole geometries we generally need a source of negative Euclidean energy. In string theory there is a natural way to obtain the required negative energy, which is to consider axion fields [9]. When a Lorentzian theory containing axions is analytically continued to Euclidean, the axion kinetic term may become negative definite which gives rise to the required negative energy-momentum tensor, see also [10–12].

In the present work [13] we will consider a consistent truncation of type IIB supergravity on S^5 down to five-dimensional maximal supergravity coupled to SO(6) gauge group [16–18]. The consistent truncation means that every solution of the five-dimensional model can be "uplifted" to a solution of full type IIB supergravity [15,19,20]. Our model will be a further (consistent) truncation to a four scalar theory coupled to AdS gravity originally introduced in [21,22] to study the holographic duals to $\mathcal{N} = 4$ Supersymmetric Yang-Mills (SYM) deformed by a mass parameter. As we will see, the model gives rise to singular domain walls, as well as regular Euclidean wormholes that have much in common with the original axionic wormholes of [9].

The embedding of axionic wormholes as AdS compactifications of 10D (or 11D) supergravity has been discussed recently in the literature [23–27]. In short, the existence of the wormhole solution relies on the existence of moduli scalars, whose metric is not necessarily positive definite. When we Wick rotate to Euclidean signature in space-time, the axions get flipped signs [28], while the other scalar fields remains untouched [12,29].

In the model studied in this paper, we will encounter a similar feature when rotating to Euclidean signature. Only one of the four scalars is a modulus that happens to be the five-dimensional dilaton. The dilaton is related to the Yang-Mills coupling constant in the dual $\mathcal{N} = 4$ SYM theory. The other scalars in our model have a nontrivial potential and play a crucial role in our construction of wormholes. Their presence makes the equations of motion much more complicated than for standard axionic wormholes. However with the help of first order BPS (Bogomol'nyi, Prasad, Sommerfield) equations that ensure supersymmetry of the solutions, we will be able to find supersymmetric Euclidean wormholes. In more detail, the solution we construct is a five-dimensional domain wall of the form

$$\mathrm{d}s_5^2 = \mathrm{d}r^2 + \mathrm{e}^{2A(r)}\mathrm{d}s_{\mathbf{R}^4}^2, \tag{1}$$

where $\mathrm{d}s_{\mathbf{R}^4}^2$ is the flat metric on $\mathbf{R}^4$ and the metric function A only depends on r. The $\mathbf{R}^4$ can be compactified to T^4 when considering the Euclidean gravity path integral. The solution we find will be described by a metric function $A(r)$ that approaches the standard asymptotic AdS form $A(r) \sim \pm r/2L$ for both $r \to \pm\infty$ connected by a region of smooth region of non-AdS space. Close to the asymptotic boundary we find a pair of (Euclidean) AdS_5 spaces dual to $\mathcal{N} = 4$ SYM. The scalar fields and the metric have asymptotic form which is consistent with masses for the three chiral fields in $\mathcal{N} = 4$ being turned on. Supersymmetry is therefore broken from $\mathcal{N} = 4$ to $\mathcal{N} = 1$. The two QFTs dual to each side of the wormhole are both $\mathcal{N} = 1^*$, but with different Yang-Mills coupling constants and different vevs turned on. A special line of solutions exists where the configuration is slightly more symmetric and the two boundary theories have the same vevs. On this line there is a very special point where the neck of the wormhole shrinks to zero size and the metric becomes singular and we recover the well-known GPPZ (Girardello, Petrini, Porrati, Zaffaroni) solution [30,31].

An important question when faced with wormhole solutions such as these ones is whether they dominate over the corresponding "disconnected geometry." Since we have been focusing on supersymmetric solutions, the disconnected geometry should also preserve supersymmetry. In our analysis we have been able to fully classify solutions to the BPS equations subject to the metric ansatz (1). It turns out that for a given set of boundary conditions which allow for a wormhole solution, there is no corresponding disconnected solution. Disconnected solutions could perhaps be found by relaxing some of the isometries built into the metric ansatz (1), but we have not carried out a general analysis. It is straightforward to check that for our BPS solutions, the regularized on-shell Euclidean action vanishes.

5D supergravity.—The supergravity model considered here is a four scalar truncation of maximal 5D supergravity with SO(6) gauge group [16–18]. The 5D SO(6) gauged maximal supergravity has been shown to arise as consistent truncation of type IIB supergravity on S^5 [15,19,20] and so any solution of the maximal supergravity can be embedded into type IIB supergravity.

The four scalar truncation discussed presently was first introduced in the holographic study of the $\mathcal{N} = 1$ mass deformation of $\mathcal{N} = 4$ SYM with all three mass parameters taken to be equal [21,22,32]. When all masses are equal the QFT possesses SO(3) flavor symmetry which (if we assume it is not spontaneously broken) can be utilized on the supergravity side to truncate the maximal theory such that the bosonic sector contains a metric and eight scalar fields [21]. A further discrete symmetry can be imposed to truncate the theory even further leaving only four scalar fields apart from the metric.

As the name suggests, the scalars of the model parametrize a four-dimensional subspace of the full 42 dimensional scalar manifold $E_{6(6)}/\mathrm{USp}(8)$. This subspace consists of two copies of the Poincaré disc which we parametrize with two complex scalar fields $z_{1,2}$ [33].

The five-dimensional supergravity action of the four scalar model takes the form

$$S = \frac{1}{16\pi G_N}\int \star\left(R - 2K_{i\bar{j}}\partial_\mu z^i\partial^\mu \bar{z}^{\bar{j}} - \mathcal{P}\right), \tag{2}$$

where the scalar potential is

$$\mathcal{P} = \frac{1}{2}\mathrm{e}^K\left(K^{i\bar{j}}D_i\mathcal{W}D_{\bar{j}}\bar{\mathcal{W}} - \frac{8}{3}|\mathcal{W}|^2\right), \tag{3}$$

with the Kähler covariant derivative defined as $D_i f = (\partial_i + \partial_i K)f$, and the Kähler metric defined by $K_{i\bar{j}} = \partial_i\partial_{\bar{j}}K$ and its inverse is $K^{i\bar{j}}$. We have written the theory in terms of the Kähler potential K and a holomorphic superpotential $\mathcal{W}$ which are given by

$$K = -\sum_{i=1}^{2}\log(2\mathrm{Im}z_i), \qquad \mathcal{W} = 3gz_2(z_1 + z_2). \tag{4}$$

The theory exhibits the scaling symmetry $z_i \mapsto \lambda z_i$ which leaves the action invariant. This is nothing but the dilatonic shift symmetry.

The maximally supersymmetric vacuum solution of the maximal five-dimensional supergravity is obtained as a critical point of this model by setting $z_1 = z_2 = i\mathrm{e}^{\varphi}$ where φ is the constant value given to the five-dimensional dilaton. For the vacuum solution, the scalar potential takes the value $\mathcal{P} = -3g^2$ and therefore the metric is AdS_5 with length scale $L = (2/g)$.

BPS equations and wormhole solutions.—We are interested in finding flat sliced supersymmetric domain wall solutions to the equations of motion of the four scalar model. To this end we assume that the five-dimensional metric takes the form (1) and assume that all scalar fields as well as the metric function A are only functions of the radial variable r. A supersymmetric solution must satisfy the first order equations [32]

$$\mathcal{E}_A \equiv A' + \frac{1}{3}W = 0, \qquad \mathcal{E}^i \equiv (z^i)' - K^{i\bar{j}}\partial_{\bar{j}}W = 0, \tag{5}$$

where the real superpotential W is defined as $W = \mathrm{e}^{K/2}|\mathcal{W}|$. We have verified that all solutions to the BPS equations are also solutions to the five-dimensional equations of motion. It should be noted at this point that in Lorentzian supergravity $\bar{z}_{\bar{\imath}}$ is the complex conjugate of z_i

and the same holds true for $\mathcal{W}$ and $\bar{\mathcal{W}}$. In this Letter we will also consider Euclidean solutions where $\bar{z}_{\bar{i}}$ is best treated as independent from z_i. In general z_i and $\bar{z}_{\bar{i}}$ still represent 2 real degrees of freedom in total. This feature has been discussed previously in, e.g., [22] but will become more apparent later when we discuss the explicit solution to the BPS equations.

In order to simplify the system of BPS equations, we introduce new field variables $z_1 = i\mathrm{e}^{\varphi+3\alpha+i\theta_1}$ and $z_2 = i\mathrm{e}^{\varphi-\alpha+i\theta_2}$. The new scalar fields have a clear interpretation from the perspective of the holographic dual field theory. In particular, φ is the five-dimensional dilaton and is dual to the marginal Yang-Mills coupling, α is dual to a scalar bilinear operator transforming in the $\mathbf{20}'$ representation of SO(6) and $\theta_{1,2}$ are dual to two fermion bilinear operators transforming in the $\mathbf{10}\oplus\overline{\mathbf{10}}$ representation.

It turns out to be useful to further define new sets of variables $t_{1,2} = \tan\theta_{1,2}$ in order to eliminate most of the trigonometric functions. Even with these new variables the BPS equations are quite lengthy and difficult to analyze. In order to make progress we replace the field α by a new variable X defined by

$$X \equiv \frac{1}{2(1+t_2^2)}\left(1 + t_1 t_2 + \sqrt{1+t_1^2}\sqrt{1+t_2^2}\cosh 4\alpha\right). \quad (6)$$

This definition of X may seem *ad hoc* at first but it is closely related to the real superpotential W. With these definitions the BPS equations take the form

$$\begin{aligned} 4\sqrt{X}(t_1') &= 3g\big(t_2 - t_1 + 2Xt_1(1+t_2^2)\big),\\ 4\sqrt{X}(t_2') &= g\big(t_1 - t_2 + 6Xt_2(1+t_2^2)\big),\\ 4\sqrt{X}(X') &= 8gX(X-1),\\ 4\sqrt{X}(A') &= -2gX(1+t_2^2). \end{aligned} \quad (7)$$

Writing the BPS equations in these coordinates has simplified them significantly enabling us to fully solve them. Note that the dilaton has been decoupled completely from the system as it does not appear on the right-hand side of any of the equations. This does not mean that the dilaton is constant, however, as its BPS equation has a complicated right-hand side.

Recall now that the scaling symmetry present in our model implies the existence of a constant of motion [34]

$$j = -\frac{g^3}{64}\mathrm{e}^{3A}(t_1 + 3t_2), \quad (8)$$

which implies that we do not have to solve explicitly the equations for both t_1 and t_2, only one combination of them suffices. For this purpose we identify another combination of the t scalars and solve for X in terms of the new variable

$$Y \equiv \frac{g^3}{64}\mathrm{e}^{3A}(t_1 - t_2), \qquad X = \frac{Y^2}{k+Y^2}, \quad (9)$$

where k is a real integration constant. We can rescale the metric function such that without loss of generality we can consider three distinct values $k = \{-1, 0, 1\}$. Next we use (9) to write

$$\frac{\mathrm{d}A}{\mathrm{d}Y} = \frac{1}{2}\frac{Y}{k+Y^2}\left(1 + 256 g^{-6}\mathrm{e}^{-6A}(j+Y)^2\right). \quad (10)$$

Finally, we remark that in the Y coordinate, the five-dimensional metric takes the form

$$\mathrm{d}s_5^2 = \frac{\mathrm{d}Y^2}{g^2(Y^2+k)} + \mathrm{e}^{2A(Y)}\mathrm{d}s^2_{\mathbf{R}^4}. \quad (11)$$

A wormhole solution is obtained if the metric function e^{6A} has two AdS_5 asymptotic regions (for $|Y|$ large) and is otherwise positive. This only happens if $k = 1$ and the metric function takes the form

$$\frac{g^6}{2^6}\mathrm{e}^{6A} = 4\left(2jY^3 - 3Y^2 - j^2 - 2 + 2a(Y^2+1)^{3/2}\right). \quad (12)$$

Since e^{6A} has at most two real roots, we have to ensure that the discriminant of the polynomial $(2jY^3 - 3Y^2 - j^2 - 2)^2 - 4a^2(Y^2+1)^3$ is negative implying it has no real roots. Combined with the condition that $a > |j|$ we find wormhole solutions if and only if $1 + j^2/2 < a$. It is now easy to see that the scalar field α is imaginary when the above condition is satisfied. In fact the condition for α being real is that $1 + j^2 \geq a^2$. The boundary of which is where the scalar α vanishes throughout and can be identified with the GPPZ solutions [30]. The two regions are completely nonoverlapping. It is interesting to note that the GPPZ solution with $j = 0$ and $a = 1$ (or $\lambda = 1$ in the notation of [32]) is infinitesimally close to being a wormhole and can be viewed as the limiting solution where the wormhole neck shrinks to zero size.

In addition to the α scalar being imaginary, the dilaton is also imaginary. In order to see that we have to integrate the BPS equation for the dilaton which for $k = 1$ takes the form

$$\varphi(Y) - \varphi_0 = \int_{-\infty}^{Y} \frac{-3i\sqrt{a^2 - j^2 - 1}\,y(y-j)}{2\sqrt{y^2+1}\left(a\sqrt{y^2+1} + jy - 1\right)\left[y^2\left(a\sqrt{y^2+1}+3\right) + a\sqrt{y^2+1} + j(y^2-3)y - 1\right]}\mathrm{d}y. \quad (13)$$

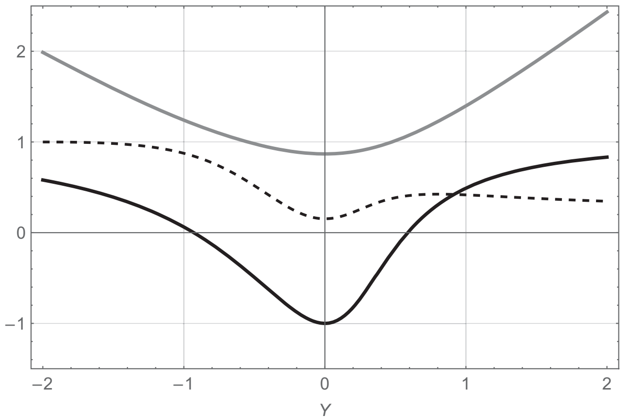

FIG. 1. A plot of a wormhole solution for $a = 2.6$ and $j = 0.8$. The function $g^2 e^{2A}/10$ is drawn in gray, $\cosh 4\alpha$ is drawn in solid black, and $\cosh 4\varphi$ in dashed black.

We have not been able to perform the integral analytically but it is easy to do numerically. We display a plot of a sample solution in Fig. 1. A generic feature of the wormhole solutions is that they are not symmetric around $Y = 0$, even for $j = 0$. This is clearly observed from Fig. 1 where the dilaton is far from being symmetric around $Y = 0$.

Field theory interpretation and on-shell action.—Recall that the scalar fields α, φ, and $\theta_{1,2}$ have a direct relation to operators in the dual field theory. More precisely if $\theta_1 = -3\phi_1 + \phi_2$ and $\theta_2 = \phi_1 + \phi_2$, and we think of $\mathcal{N} = 4$ SYM in $\mathcal{N} = 1$ language, then ϕ_2 is dual to the gaugino bilinear and ϕ_1 is dual to the three chiral multiplet fermion bilinear. The latter are all equal since we assumed that the SO(3) flavor symmetry preserved by the equal mass $\mathcal{N} = 1^*$ Lagrangian is not spontaneously broken. By performing a expansion of the fields around the AdS_5 asymptotics we can identify the sources and vevs given to the dual operators for our solution. A general UV expansion compatible with the BPS equations takes the form

$$A = -\frac{1}{2}\log\epsilon + \mathcal{O}(\epsilon^2), \qquad \phi_1 = m\epsilon^{1/2} + \mathcal{O}(\epsilon^2),$$
$$\phi_2 = w\epsilon^{3/2} + \mathcal{O}(\epsilon^2), \qquad \alpha = v\epsilon + \mathcal{O}(\epsilon^2), \tag{14}$$

and φ is constant at leading order. Here ϵ is the small parameter controlling the distance from the asymptotic boundary. From this expansion we see that the parameter m is the mass given to the chiral fields whereas w is the gaugino vev and v is the so-called chiral condensate, i.e., the vev of the scalar bilinear in the chiral multiplets.

For the wormhole solutions with $k = 1$, we have two asymptotic regions that are located at $Y \to \pm\infty$. Expanding our solution we find the dimensionless quantities

$$\frac{w_\pm}{2m^3} = (j^2 + 1) \pm aj, \qquad \frac{v^2}{m^4} = (j^2 + 1) - a^2, \tag{15}$$

where we have denoted the two gaugino condensates that are encountered in the two asymptotic regions $Y \to \pm\infty$ by $w_\pm$. The dimensionless chiral condensate v/m^2, is the same in both regions. Regular wormhole solution exist only when the chiral condensate v/m^2 is imaginary. The conclusion is that the Euclidean wormhole solution is a bulk geometry that connects two copies of mass-deformed $\mathcal{N} = 4$ SYM where some of the boundary conditions (including the Yang-Mills coupling constant) are different.

An important aspect for the evaluation of the gravitational partition function, and of the free energy of the dual theory, is the on-shell action of our wormholes solutions. Adding the Gibbons-Hawking term to the action and performing a partial integration, the action can be rewritten in terms of the squared BPS equations

$$\mathcal{L} + \mathcal{L}_{\mathrm{GH}} = -e^{4A}\left[12\mathcal{E}_A^2 - 2K_{i\bar{j}}\mathcal{E}^i\bar{\mathcal{E}}^{\bar{j}}\right] - 2\partial_r(e^{4A}W),$$

where W is the real superpotential we defined before $W = e^{K/2}|\mathcal{W}|$. Evaluating this on-shell, the BPS equations set to zero the first two terms, leaving only the total derivative. The total derivative term leads to a divergent expression which must be regulated. As explained in [22,36,37] the correct supersymmetric counterterm (when the holographic boundary is flat) should be chosen to exactly cancel the total derivative. This implies that for all supersymmetric regular wormhole solutions found in this Letter, the on-shell action vanishes.

Final comments.—As anticipated in the introduction, our wormhole solutions are supported by a negative term in the energy-momentum tensor. In fact, the dilaton φ and the field α are imaginary while preserving the reality of the metric as well as the action. This is interpreted as a Wick rotation on target space which must be simultaneously performed when Wick rotating space-time. The target space metric exhibits a pair of translation symmetries for α and φ, which therefore appear as axions. However it must be noted that the scalar potential depends nontrivially on α and so the shift "symmetry" of α is not a true symmetry of the theory. Nevertheless, according to the prescription in [24,26], the Wick rotation of Lorentzian supergravity to Euclidean should be accompanied with a similar Wick rotation in target space $\alpha \to i\alpha$ and $\varphi \to i\varphi$ [38]. The Wick rotation affects the potential but is still real.

While the space-time signature becomes Euclidean, the target space signature is now Lorentzian. This is not particularly surprising if we remember that in Euclidean signature, the R symmetry of $\mathcal{N} = 4$ SYM is not SO(6) but rather SO(1,5). This is also consistent with the fact that the holographic dual to Euclidean $\mathcal{N} = 4$ SYM is described by so-called type IIB* supergravity [39–41]. This theory has a metric with space-time signature (9,1) but some of the form fields have negative kinetic terms (and can therefore be thought of as being analytically continued from standard type IIB supergravity). In particular the ten-dimensional axion-dilaton parametrize the coset space $SU(1,1)/SO(1,1)$. It is a subject of future work to uplift the wormholes to ten dimensions using the results of [15,19,20]. The fact that the five-dimensional dilaton is imaginary may appear worrisome when interpreted in ten dimensions. Even though the relation between the five-dimensional dilaton and the ten-dimensional one is rather complicated, asymptotically they are simply related. It would therefore appear that the ten-dimensional dilaton is imaginary for the wormhole solutions which is troubling. It turns out there is a simple remedy for this problem by employing an $SL(2,\mathbf{R})$ transformation that renders the dilaton real but turns on an imaginary axion consistent with being a solution of type IIB* [35].

Since our solutions preserve supersymmetry and are regular we do not expect any instabilities to arise and question the validity of them. By all accounts they should then contribute to the Euclidean path integral. Since we did not find disconnected geometries with the same boundary conditions as the wormholes, we are unable to answer whether the wormholes dominate or not. If they dominate, then it leads to the well-known factorization puzzle in holography [7,8] in this case for deformations of AdS_5 dual to the $\mathcal{N}=1^*$ theory. How this puzzle is resolved is a open question at this stage. One possibility is that fermion zero modes in the spectrum cause the wormhole contribution to the path integral to vanish. This was indeed observed in [42] and in that case it was related to supersymmetry being broken.

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