Meromorphic Functions and the Topology of Giant Gravitons

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Abstract:
Using Mikhailov’s map from holomorphic functions to supersymmetric D3-brane solutions, we show how to construct giant gravitons in $\text{AdS}_5 \times S^5$ with toroidal topologies. In the $\frac{1}{4}$-BPS sector we show that these are always of the form $\sharp^K (S^2 \times S^1)$, and in the limit in which this becomes a set of $m + n$ perpendicular spherical giants re-connected near to their intersections, we find $K$ in terms of $m, n$. In the $\frac{1}{8}$-BPS sector we find a similar class of solutions.

1 Introduction

The best understood sector of AdS/CFT [1] concerns closed fundamental strings, which in the gauge theory are single-trace operators of length much less than $N$ [2]. D-branes are heavier objects corresponding to operators of length order $N$, the most studied of which are spherical branes known as giant gravitons [3], dual to determinant operators and generalisations known as Schur polynomials [4].

The worldsheet of any closed string state is a cylinder $\mathbb{R} \times S^1$, but for D3-branes any $\mathbb{R} \times \mathcal{M}$ is in principle possible, with $\mathcal{M}$ a closed 3-manifold. Such manifolds can have considerably more complicated topology than closed 1-manifolds, and it would be fascinating to understand the emergence of topology from the dual SYM operators. Motivated by this, the goal of this letter is to explore what topologies occur in giants with a given amount of supersymmetry. We are particularly interested in solutions created by a localised modification of a set of intersecting spherical giant gravitons, as this seems the most tractable limit.

We begin by recalling the map given by Mikhailov [5]: Any analytic function $f : \mathbb{C}^3 \to \mathbb{C}$ defines a supersymmetric D3-brane solution in $\mathbb{R} \times S^5 \subset \text{AdS}_5 \times S^5$ as the surface

$$f(e^{-it}Z_1, e^{-it}Z_2, e^{-it}Z_3) = 0, \quad \sum_i |Z_i|^2 = 1$$

(1)
where $Z_i = r_i e^{i \phi_i}$ are the 3 complex embedding co-ordinates for $S^5$. The degree of supersymmetry is given by the number of arguments: $f(Z_1)$ gives a $\frac{1}{2}$-BPS solution, $f(Z_1, Z_2)$ $\frac{1}{4}$-BPS, and $f(Z_1, Z_2, Z_3)$ gives an $\frac{1}{8}$-BPS solution.

The usual sphere giant graviton is given in this language by

$$f(Z_1) = Z_1 - \alpha.$$  \hfill (2)

This constrains $Z_1$ completely, so the worldvolume (at a given time) is the $S^3$ parameterised by $Z_2$ and $Z_3$ subject to $|Z_2|^2 + |Z_3|^2 = 1 - \alpha^2$. The only time-evolution is that it rotates in the $Z_1$ plane; the maximal giant graviton has $\alpha = 0$ and is thus stationary, while in the opposite limit $\alpha \to 1$ the brane collapses down to a point particle on a lightlike trajectory. A function $f(Z_1)$ with several zeros will lead to a number of concentric spherical giants.

We refer to (2) as the case $(1, 0, 0)$: one $Z_1$ giant. The next section studies the effect of adding to this terms depending on $Z_2$, and then takes a limit in which these give $n$ intersecting $Z_2$ giants: cases $(1, n, 0)$. After that we consider arbitrarily many intersecting $Z_1$ and $Z_2$ giants, cases $(m, n, 0)$ (section 3), and finally the addition also of $Z_3$ giants (section 4). We give a concise statement of our results in section 5.

2 Quarter-BPS Class $(1, n, 0)$

To begin constructing topologically nontrivial solutions using Mikhailov’s method, in this section we add to the spherical giant’s $f(Z_1)$ a meromorphic function of $Z_2$. Consider first the function

$$f(Z_1, Z_2) = Z_1 - \alpha + \frac{\epsilon}{Z_2}.$$  \hfill (3)

We may assume $\alpha, \epsilon > 0$ and, since the motion of the brane is rigid, we need only discuss its topology at time $t = 0$.

Let us parameterise the D3-brane worldvolume by the $\phi_3$ circle and some portion of the $Z_2$ plane. (Solving $f = 0$ fixes $Z_1$ in terms of $Z_2$, and $\sum r_i^2 = 1$ fixes $r_3$.) We can read off the topology of the brane from the topology of the area of the $Z_2$ plane thus covered: let $\Sigma$ be the area where $r_3 \geq 0$. For example, the spherical giant graviton (2) clearly has for $\Sigma$ the disk $|Z_2| \leq 1 - \alpha^2$.

The effect of turning on the pole is to make a hole in the base space $\Sigma$, thus increasing its genus (see figure 1). This may be understood by saying that in a neighbourhood of the pole, the term $\epsilon/Z_2$ is so large that there are no solutions $|Z_1| \leq 1$. Notice immediately that this means that the pole itself is not on the worldvolume of the brane.
Figure 2: Plots showing $\Sigma$ with one quintuple pole, (5) with $N = 5$. Increasing $\epsilon$ we pass from the torus on the left to five spheres on the right; the middle picture is close to $\epsilon_{\text{crit}}$. Parameters are $\alpha = \frac{1}{2}$ and $\epsilon = \frac{1}{1000}, 0.03833, 1.5$.

Figure 3: Plots showing $\Sigma$ for the case $(1, 3, 0)$, using (6) with three poles at $\beta_1 = -\frac{1}{2}$, $\beta_2 = 0$ and $\beta_3 = \frac{1}{2}$. The residues are $\epsilon, \epsilon, -\epsilon$ respectively, with $\epsilon$ increasing from $\frac{1}{12}$ (left, giving $M = \mathbb{Z}^3(\mathbb{S}^2 \times \mathbb{S}^1)$) to $\frac{1}{3}$ (centre, giving $\mathbb{S}^2 \times \mathbb{S}^1$) to $\frac{1}{3}$ (right, $\sqcup_3 \mathbb{S}^3$), and $\alpha = \frac{1}{2}$. Notice that the holes in $\Sigma$ formed by residues of opposite signs attract, while those of the same sign repel. The lines $C_3, C_4$ in $\Sigma$ each lift to a separating $S^2$ in $M$.

To analyse this more carefully, it is easy to show using (3) that $\Sigma$ is given by

$$r_2^4 + r_2^2(\alpha^2 - 1) + \epsilon^2 \leq 2c_\alpha r_2 \cos \phi_2.$$  \hspace{1cm} (4)

Drawing graphs of the left- and right-hand sides in terms of $r_2^2$, when $\epsilon = 0$ certainly there are two intersections. Increasing $\epsilon$, there is a range $0 < \epsilon < \epsilon_{\text{crit}}$ in which there are two intersections $r_2 > 0$ for all $\phi_2$, followed by a range $\epsilon_{\text{crit}} < \epsilon < \epsilon_{\text{max}}$ in which there are two intersections at $\phi_2 = 0$ but not at $\phi_2 = \pi$. For larger $\epsilon$ there are no intersections. This progression is shown in figure 1.

For $\epsilon < \epsilon_{\text{crit}}$ the topology of the brane is $S^2 \times S^1$. The incontractible cycle $C_1$ (more or less the $\phi_2$ circle) is the $S^1$ factor, while a radial line in $\Sigma$ gives the $S^2$ factor — this is an interval fibered with the $\phi_3$ circle, which shrinks to zero at either end. It is marked $C_2$ in the figure.

The simplest generalisation is to consider a higher-order pole:

$$f(Z_1, Z_2) = Z_1 - \alpha + \frac{\epsilon}{(Z_2)^N}.$$  \hspace{1cm} (5)

This leads to the same topology as the single pole, for small $\epsilon$, but the geometry has a symmetry $Z_2 \rightarrow e^{2\pi i/N} Z_2$. Because of this, in the regime $\epsilon_{\text{crit}} < \epsilon < \epsilon_{\text{max}}$ there will be $N$ separate (deformed) 3-spheres. Figure 2 shows the case of $N = 5$.

We can also consider several poles:

$$f(Z_1, Z_2) = Z_1 - \alpha + \sum_{j=1}^{n} \frac{\epsilon_j}{Z_2 - \beta_j}.$$  \hspace{1cm} (6)

For small enough $\epsilon_j$ the analysis very close to each pole will be similar to that for one pole: expanding in $r_{2\beta} = |Z_2 - \beta_j|$ will give us (4) plus terms higher order in $r_{2\beta}$. Thus for any set of
n poles (located at $\beta_j$ such that $\alpha^2 + |\beta_j|^2 < 1$) there exist residues $\epsilon_j \neq 0$ such that $\Sigma$ is a disk with $n$ holes. Cutting $\Sigma$ along lines such as $C_3, C_4$ in figure 3, so that each hole is isolated, we see that the resulting topology is a connected sum

$$\mathcal{M} = \sharp^n(S^2 \times S^1).$$

Note that all poles are outside $\Sigma$, so everywhere on the worldvolume the function $f$ is analytic.

With multiple poles the progression as we increase $\epsilon$ can be quite complicated, and can produce several disconnected pieces. The case of poles at $Z_2 = 0, \pm \frac{1}{2}$ is shown in figure 3. Notice that the holes created by $\epsilon_2/Z_2$ and $\epsilon_3/(Z_2 - \beta_3)$ merge with each other in the middle picture. The same effect can be produced by moving them together at fixed $\epsilon$: when $\beta_3 \to 0$ these two approach (5). We study this kind of degeneration limit extensively below.

Returning for a moment to our simplest case (3), there are two more distinct degeneration limits given by the two cycles shown:

- As $\epsilon \to \epsilon_{\text{crit}}$, $C_2$ becomes small in a throat which is locally $S^2 \times \mathbb{R}$. This is the right geometry to be interpreted as the effect of some strings pinching the brane into a torus. In the case $\alpha = 0$ the brane becomes an equally thin tube all around, approaching the circular spinning string solution

$$Z_1 = \frac{1}{\sqrt{2}} e^{i(t+\sigma)}, \quad Z_2 = \frac{1}{\sqrt{2}} e^{i(t-\sigma)}, \quad Z_3 = 0.$$  

The toroidal brane may thus be thought of as an embiggened circular string. But note that all D3-branes described by (1) carry no worldsheet electric field $F_{01}$, and thus the string solution here is not an F-string.

- As $\epsilon \to 0$, instead $C_1$ becomes small in a throat locally $S^1 \times \mathbb{R}^2$. Here it is useful to think of the $\epsilon \neq 0$ function (3) not as the addition of a meromorphic term to (2) but (multiplying through by the denominator) as a the addition of a small term to a factorised polynomial. That is, $f = Z_1Z_2 + \epsilon$ describes exactly the same D3-brane as (3), but in the limit $\epsilon \to 0$ more obviously approaches $f = Z_1Z_2$, which is a pair of intersecting maximal sphere giants.

The effect of infinitesimal $\epsilon$ is localised near to their intersection: At $Z_2 \neq 0$ it changes $Z_1 = 0$ to $Z_1 = \epsilon/Z_2$, perturbing the $Z_1$ giant smoothly away from maximality (and likewise $Z_2$). But the effect near to $Z_1 = Z_2 = 0$ is not smooth, as $\cup_2 S^3$ is re-connected so as to give topologically $S^2 \times S^1$.

For three poles in (6) (and taking the numbers used in figure 3) the view suggested by the limit $\epsilon \to 0$ is of four intersecting branes, and hence we refer to this as the case $(1, 3, 0)$:

$$f(Z_1, Z_2) = (Z_1 - \frac{1}{2})(Z_2 + \frac{1}{2})Z_2(Z_2 - \frac{1}{2}) + \epsilon (Z_2^2 - Z_2 - 1) 
\approx (Z_1 - \frac{1}{2})Z_2(Z_2^2 - \frac{1}{4}) + \epsilon'$$

The three $Z_2$ branes intersect the $Z_1$ brane at different places, and since $\epsilon \neq 0$ modifies the solution appreciably only near to the intersection, it is natural that the effect of several $Z_2$ branes is very simply related to the effect of one, (14). In the next section we study more general cases in this limit, allowing also several $Z_1$ branes.

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1 Recall that the notion of a connected sum is this: If cutting a 3-manifold $M$ along an $S^2$ separates the manifold into $M_1' \cup M_2'$, and $M_1$ is $M_1'$ with a 3-ball glued to its boundary, then we write $M = M_1 \sharp M_2$. The sphere is the identity in the sense $M = M \sharp S^3$. Every (oriented closed connected) 3-manifold has a unique decomposition as a sum of prime manifolds, primeness meaning that every separating $S^2$ bounds a ball.

The connected sum of 2-manifolds, which we write $\natural$, is defined by similarly cutting along $S^1$. This gives rise to the genus classification of surfaces $S^2, T^2, T^2 \natural T^2, \cdots$.
Figure 4: Branch cut for the $(3,1,0)$ case (9). Each sheet of $\Sigma$ on the left can be turned inside-out to give a wedge as shown. (The numbers label boundary components.) Glueing these back together, the result is a disk with three holes, drawn schematically at the right, and equivalent to figure 3a. (Note that the cuts labelled $A = A'$ etc. are not the $S^2$ glue lines of the connected sum.)

3 Class $(m, n, 0)$

Figure 3 above shows $\Sigma$ for the case $(1,3,0)$. Let us now analyse the case $(3,1,0)$, which must be equivalent. The simplest example is

$$f = (Z_1^3 - \alpha^3)Z_2 + \epsilon. \quad (9)$$

Solving for $Z_1 = \sqrt[3]{\alpha^3 - \epsilon/Z_2}$, if we again call the area of the $Z_2$ plane occupied by the solution $\Sigma$, this is now a Riemann surface with three sheets. Each sheet is a disk with one hole, and the branch cut (from $Z_2 = 0$ to $\epsilon/\alpha^3$) runs from a point in the hole to a point inside $\Sigma$. Drawing the connections as in figure 4, it is clear that $\Sigma$ has three holes, thus we recover $M = \#_3(S^2 \times S^1)$ as desired. The same procedure works equally well for branch cuts of any order.

Before moving on to a new case, we make the following observation: The connected sum of two 3-manifolds $M, N$ can be regarded as the effect of connecting a tube $S^2 \times I$ (i.e. $S^3$ with two punctures) between any point on one and any point on the other:

$$M \# N = M \# S^3 \# N.$$  

Connecting a similar tube between two points on the same manifold $M$ instead has the effect of adding one term $S^2 \times S^1$:

$$M + (S^2 \times I \text{ handle}) = M \# (S^2 \times S^1). \quad (10)$$

In the notation of the figures, where $\square = S^2 \times S^1$ and $\square = S^2 \times I \text{ tube}$, this reads

$$\square \square \square \square \square = \square \square \square \square \square = \#_2(\square).$$

Now consider the case $(2,2,0)$, starting with

$$f = (Z_1^2 - \alpha^2)(Z_2^2 - \beta^2) + \epsilon. \quad (11)$$

Clearly $\Sigma$ has two sheets, each with two holes, connected by a pair of branch cuts. To analyse this, we can split each sheet along surfaces $C_i$ like those used in the previous section — see figure 5. This gives two copies of the $(2,1,0)$ case, $\#^2(\square)$, connected in two places. Re-connecting along $C_3$ on the upper sheet gives the connected sum of the two pieces, and re-connecting the lower sheet adds a handle of the type just discussed. In all we get

$$M = \left[\#^2(S^2 \times S^1)\right] \# \left[\#^2(S^2 \times S^1)\right] + (S^2 \times I \text{ handle}) = \#^5(S^2 \times S^1).$$
Figure 5: Branch cuts for case (2, 2, 0). The first plot shows two square-root cuts each of which can be treated as in figure 4, giving the vertical connections; the horizontal connections are from the glue line $C_3$. The second plot shows the same branch points connected the other way. In this case we can pull the lower sheet of $\Sigma$ through the cut to obtain the figure on the right. (This happens within the dashed line. The circular boundary component was the outer boundary of the lower sheet.) Now $\Sigma$ is a disk with three holes plus two handles, giving the same topology.

Figure 5 shows this procedure. Instead of (11) it uses $f = (Z_2^2 - \alpha^2) + \epsilon/(Z_2 - \beta) + i\epsilon/(Z_2 + \beta)$, with $\alpha = \frac{1}{2}$, $\beta = \frac{1}{4}$ and $\epsilon = \frac{1}{9}$, to have a convenient arrangement of branch points. It also shows an alternative argument to check that the choice of how we draw the branch cuts does not matter.

Generalising to the case $(m, n, 0)$, the topology is

$$\mathcal{M} = \sharp^K(S^2 \times S^1), \quad K = mn + (m - 1)(n - 1). \quad (12)$$

The counting comes from drawing a grid of $\bigcirc$ and connecting horizontally (as in figure 3) and vertically (as figure 4). So far we have assumed that the $m + n$ intersecting branes are all at distinct positions, or in other words we considered only single poles. In the last section, allowing instead higher-order poles (5) did not change the topology, but this is no longer true here. We can investigate this by moving poles to co-incide. There are two ways to do this in (11), taking either $\alpha \to 0$ or $\beta \to 0$, and these must be equivalent. Solving for $Z_1$, the branch points are located at

$$Z_2 = \pm \beta, \quad \pm \sqrt{\beta^2 + \epsilon/\alpha^2}.$$  

For small $\epsilon$ the second pair are inside $\Sigma$, giving the analysis above. But (holding $\epsilon$ fixed and) taking the limit $\alpha \to 0$, these move off to infinity, giving cuts all the way across $\Sigma$. In the limit $\beta \to 0$, instead the holes in $\Sigma$ merge into one (as happened in figure 3). Both situations are drawn in figure 6, and each leads to $\mathcal{M} = \sharp^3(S^2 \times S^1)$.

It is now clear how to treat any set of poles of any order. Write $n$ for the number of separated poles in $Z_2$, and $N$ for their total order. (Two double poles thus give $N = 4$, $n = 2$, as does one single and one triple pole.) Similarly write $M \geq m$ for poles in $Z_1$. Drawing an $m \times n$ grid of $\bigcirc$, the number of vertical connections is $M$, and horizontal $N$ — see figure 7. Then counting the holes we get

$$\mathcal{M} = \sharp^K(S^2 \times S^1), \quad K = 1 + M(n - 1) + N(m - 1). \quad (16)$$

This change from (12) is a result of holes in $\Sigma$ merging with each other. We learned in section 2 that the effect of increasing $\epsilon$ is similar. Thus we expect that, for a completely general $\frac{1}{4}$-BPS giant, the topology will still be $\sharp^K(S^2 \times S^1)$ for some $K$.  


Figure 6: Degeneration limits of the $(2, 2, 0)$ case (11), drawing always the upper sheet of $\Sigma$. For the central picture (small $\beta$) we have two disks connected by 4 handles (labelled $A \ldots D$), while for the right-hand picture (small $\alpha$) we have 2 tori connected by 2 handles. (The initial picture is $\alpha = \beta = \frac{1}{2}, \delta = \frac{1}{11}$, and for each limit drawn “small” means $\frac{1}{3}$.)

$M = m = N = n = 3, K = 13$

$N = 3, n = 2, K = 10$

$M = N = 3, m = n = 2, K = 7$

Figure 7: Degeneration of the case $(3, 3, 0)$. Starting with three distinct poles at $Z_1 = \alpha_i$ and three at $Z_2 = \beta_j$, we first allow two $\beta$ poles to merge into a double pole, and then two $\alpha$ poles likewise. $\mathcal{M}$ is given by (16) with the numbers shown.

4 Eighth-BPS

We now turn wish to turn on at least one intersecting $Z_3$ giant, and take a similar small-$\epsilon$ limit. Using what we have learned, we can immediately treat all cases $(1, n, 1)$ together. Consider

$$f = Z_1 Z_3 + \epsilon h(Z_2)$$

where $h$ is a function with $n$ poles. Clearly $f = 0$ fixes $\phi_+ = \phi_1 + \phi_3$ and the product $r_1 r_3$ in terms of $Z_2$, while $\phi_- = \phi_1 - \phi_3$ is unconstrained. We can solve for $r_1$ and $r_3$ as separate functions of $Z_2$ by writing $\sum_i r_i^2 = 1$ as

$$(r_1 \pm r_3)^2 = 1 - r_2^2 \pm 2 r_1 r_3 \equiv H_\pm$$

and substituting in $r_1 r_3 = \epsilon |h|$. This gives

$$(r_1, r_3) = \frac{1}{2} \left( \sqrt{H_+} \pm \sqrt{H_-}, \sqrt{H_+} \mp \sqrt{H_-} \right).$$

At a point $Z_2$ for which $H_- \geq 0$ there are two solutions (coalescing when $H_- = 0$), while for $H_- < 0$ there are none. Notice that where $h$ has a pole, $H_- \to -\infty$, and thus the neighbourhood of such points will be excluded. (And for small $\epsilon$, each such hole will be small.) Define $\Sigma$ to be two copies of the area of the $Z_2$ plane for which $H_- \geq 0$, sewn up along the boundary. For one pole this is a simple torus, while for $n$ poles $\Sigma = \frac{2^n}{3} T^2$.

In the $\frac{1}{2}$-BPS case we always had an $S^1$ fibred over $\Sigma$, shrinking to a point on $\partial \Sigma$. Fitting with the fact that this boundary is now empty, the $\phi_-$ circle here never shrinks to a point. To
check, note that the metric is

\[ ds^2 = \sum_{i=1}^{3} (dr_i^2 + r_i^2 d\phi_i^2) = 2(r_1^2 + r_3^2) d\phi_2^2 + \ldots. \]

Thus for the length of the \( \phi^- \) circle to be zero we need \( r_1 = r_3 = 0 \), which implies \( r_2 = 1 \), and this is never part of \( \Sigma \). We conclude that the topology is

\[ \mathcal{M} = (\mathbb{Z}^n T^2) \times S^1. \]  

(14)

Note that all such three-manifolds are prime — the connected sum here is the two-dimensional one, and to separate the manifold (nontrivially) we must cut along \( T^2 \) surfaces.

While the area of the \( \mathbb{Z}_2 \) plane involved here is not identical to that for the \( \frac{1}{4} \)-BPS case of section 2, the spirit is clearly very similar: each pole increases the genus of the base space \( \Sigma \). It is natural to ask how much of our analysis of section 3 still holds. To understand this we begin by re-analysing \( f = Z_1 Z_2^2 - \beta^2 \) \( Z_3 + \epsilon \), the case \( (1, 2, 1) \). Solving for \( Z_2 \) we get

\[ Z_2 = \pm \beta \sqrt{Z_1 Z_3 - \epsilon / \beta^2} / \sqrt{Z_1 Z_3}. \]

It is natural to think of this as having a branch cut in the \( \mathbb{Z}_4 = \mathbb{Z}_1 \mathbb{Z}_3 \) plane. Fixing our position on this plane fixes \( Z_2 \) up to a choice of sheets, after which we still have (at a generic point) two solutions \( (r_1, r_3) \). Figure 8 shows the upper-sheet part of \( \Sigma \) which, when glued along the branch cut drawn, gives a double torus, and thus the same topology as before. The angle in the \( S^1 \) factor is still \( \phi_1 - \phi_3 \).

One extension beyond (13) is now fairly obvious. If we consider\(^2\)

\[ f = 1 + \sum_{k=1}^{m} \frac{\epsilon_k'}{Z_1 Z_3 - \gamma_k} \sum_{j=1}^{n} \frac{\epsilon_j'}{Z_2 - \beta_j}. \]

then we can re-use all the \( \frac{1}{4} \)-BPS analysis. Just as we replaced \( \Sigma \) of figure 3 with two copies glued along their edges to get (14), similarly replace \( \Sigma \) of figures 5, 7 with their closed cousins. We get \( (\mathbb{Z}_n T^2) \times S^1 \) with the same \( K \) as before.

But it is essential here that \( f \) contains only the product \( Z_1 Z_3 \). The more natural class

\(^2\) Note aside that we could likewise consider \( \frac{1}{4} \)-BPS solutions of the form \( f = (Z_1 - \alpha)(Z_4 - \gamma) + \epsilon' \). For small \( \gamma \) this can give a double torus \( \sharp^2(S^2 \times S^1) \), but small \( \epsilon' \) here does not guarantee that \( \epsilon \) in (18) is small.
\( (m, n, 1) \) of solutions

\[
f = Z_3 + \sum_{i=1}^{m} \frac{\epsilon_i'}{Z_1 - \alpha_i} + \sum_{j=1}^{n} \epsilon_j \frac{Z_2 - \beta_j}{Z_2 - \beta_j}
\]

will break the \( S^1 \) symmetry in (14). We leave the analysis of this, and completely general \( \frac{1}{8} \)-BPS cases, for future work.

5 Conclusion

The main result of this letter is as follows:

Let \( g(Z_1, Z_2) \) be a meromorphic function with \( m \) distinct poles at \( Z_1 = \alpha_i \), and write \( M \) for the number of poles counting multiplicity. Similarly let \( n \) and \( N \) count the poles at \( Z_2 = \beta_j \). We require \( m, n \geq 1 \), and \( |\alpha_i|^2 + |\beta_j|^2 \leq 1 \) \( \forall i, j \).

Consider the \( \frac{1}{4} \)-BPS giant described by

\[
f(Z_1, Z_2) = 1 + \epsilon g(Z_1, Z_2).
\]

(15)

For sufficiently small \( \epsilon \), this has topology specified by the prime decomposition

\[
\mathcal{M} = \sharp^K (S^2 \times S^1), \quad K = 1 + M(n - 1) + N(m - 1).
\]

(16)

As \( \epsilon \) is increased, generically \( K \) will decrease, and the brane may break up into several disjoint pieces. All pieces are either spheres or connected sums of \( (S^2 \times S^1) \):

\[
\mathcal{M} = \bigsqcup_i \sharp^{K_i} (S^2 \times S^1) \bigsqcup_j S^3.
\]

(17)

We found it convenient to deal with a function \( f \) with poles, which in some sense repel the base space \( \Sigma \) thus creating holes in the brane. The same solutions can equivalently be specified by polynomial functions of the form

\[
f(Z_1, Z_2) = \prod_{i=1}^{m} (Z_1 - \alpha_i)^{\mu_i} \prod_{j=1}^{n} (Z_2 - \beta_j)^{\nu_j} + \epsilon \text{poly}(Z_1, Z_2).
\]

(18)

Clearly \( \epsilon = 0 \) gives a factorised \( f \) and thus a set of intersecting spherical giants (2). The effect of small \( \epsilon \) is to suppress all but the simplest kind of interactions: the topology is unchanged when the last term here is replaced by a constant. Nevertheless what we have observed is that the effect of increasing \( \epsilon \) is quite simple: the tori degenerate (reducing \( K \)) and ultimately split into disjoint spheres. Thus we believe that (17) applies to generic polynomial functions \( f(Z_1, Z_2) \).

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3 We could weaken this condition to allow for cases where not every pair of branes intersect in the limit \( \epsilon \to 0 \); this has the effect of deleting some nodes from the corners of the lattice shown in figure 7, and thus reducing \( K \), but not otherwise changing the topology.

4 This is true if the residues of \( g \) are constants, in which case \( \sum_i K_i \leq K \) and \( L + L' \leq MN \) in (17). But if the numerator of \( g \) is of sufficiently high order then \( K \) may increase.

5 The poles are thus never on the worldvolume, so \( f \) is locally analytic, which is enough to guarantee a solution to the equations of motion from (1).

6 Here \( M = \sum_i \mu_i \) and \( N = \sum_j \nu_j \).
For $\frac{1}{8}$-BPS geometries we have more limited results. The generalisation which can be treated by borrowing much of the analysis from above is

$$f(Z_1, Z_2, Z_3) = 1 + \epsilon g(Z_4, Z_2), \quad Z_4 = Z_1 Z_3$$

with $g$ defined as in (15). For small enough $\epsilon$ the resulting topology is

$$\mathcal{M} = [\mathbb{Z}^K(S^1 \times S^1)] \times S^1$$

where $K$ is as in (16). Notice that none of these topologies can occur in the $\frac{1}{4}$-BPS case. Generalising this to allow other combinations of $Z_3$ branes (such as (18) with $\prod_k (Z_3 - \gamma_k)$ inserted) is an open problem. But it seems clear that the form of $\mathcal{M}$ will change, and in particular will not have an overall $S^1$ factor.

While our focus in this letter has been entirely on the classical membranes described by (1), a detailed quantisation of the moduli space of Mikhailov solutions was carried out in [6], and used to draw conclusions about the spectrum of $\frac{1}{8}$-BPS states in $\mathcal{N} = 4$ SYM. It would be of great interest to pursue the relationship between these results and ours.

This work forms part of a larger research program aimed at understanding how local and global properties of spacetime are encoded in gauge theory. For recent work in this direction, see [7] and [8] and references therein. In this context, it would be nice to see how the topologies studied here emerge in operators dual to Mikhailov’s giants in SYM. There too, perhaps our small-$\epsilon$ limit is likely to be the tractable one.

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7 From the topologies written down here it is trivial to obtain the homology groups. The Betti numbers are:

- $\frac{1}{4}$-BPS: $b_0 = b_3 = 1, \quad b_1 = b_2 = K$
- $\frac{1}{8}$-BPS: $b_0 = b_1 = 1, \quad b_1 = b_2 = 2K + 1$

For the case $(1, 1, 0)$, the generators of $H_1$ and $H_2$ are cycles $C_1$ and $C_2$ in figure 1. It is easy to draw similar cycles in figure 3’s case $(1, 3, 0)$. Similar cycles drawn in figure 8’s case $(1, 2, 1)$ will all be 1-cycles.
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