(Euclidean) Wormholes and Wilson Loops

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and in progress

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Why to study (Euclidean) Wormholes

Wormholes are interesting (exotic) solutions of GR + matter

- Proposed physical effects due to wormholes
  - They lead to a non-trivial topology of space(time)
  - Implications on the information paradox? - Connect the black hole interior with exterior?
  - Affecting the low energy coupling constants? (α-parameters) [Lavrelashvili, Rubakov, Tinyakov, Coleman, Hawking ...] - Resolving the Cosmological Constant problem?
  - Related to cosmological spacetimes upon analytic continuation

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    - Resolving the Cosmological Constant problem?
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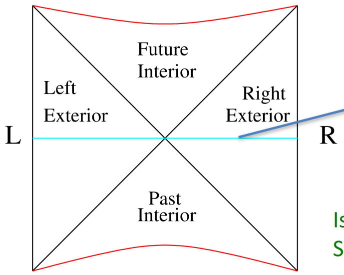
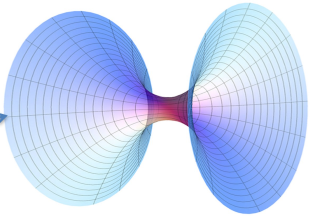
- Different types of wormholes
  - Lorentzian vs Euclidean
  - Macroscopic multi-boundary geometries (saddles) vs.
    Microscopic "gas of wormholes" (α-parameters)
  - Different characteristic scales
    $L_P \ll L_W \sim L_{macro}$ vs. $L_P \leq L_W \ll L_{macro}$ ex: $L_{macro} = L_{AdS}$

- Our main focus will be macroscopic (Euclidean) wormholes in the context of holography (AdS/CFT)

Lorentzian wormholes or "ER = EPR"

- Einstein - Rosen Bridge: Connects the two sides of the eternal black hole





Israel Maldacena
Susskind

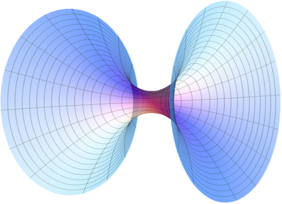
- Wormhole = Einstein Podolski Rosen pair of two black holes in a particular entangled state of two non-interacting QFT's:

$$|\Psi\rangle = \sum_n e^{-\beta E_n/2} |E_n\rangle_L^{CPT} \times |E_n\rangle_R$$

- Large amounts of entanglement can give rise

to a geometric connection!

Euclidean Wormholes (saddles)



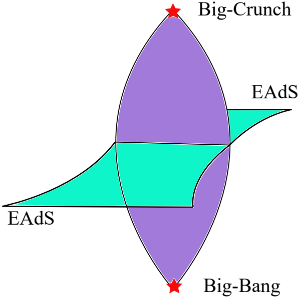
- Now the complete space (Euclidean) is a wormhole
  No additional Lorentzian time direction
- To have such solutions, one needs locally negative Euclidean Energy to support the throat from collapsing
- Such energy can be provided by axionic fields or "magnetic" fluxes

- Existence of solutions in various dimensions/setups and with different asymptotics (i.e. flat vs. AdS)
- Some can even be embedded in the standard model + gravity
- A subset of those is perturbatively stable
  [Marolf-Santos, Hertog-Van Riet ...]
- What about their analytic continuation into Lorentzian?
  This gives a further reason why Euclidean wormholes are interesting (see [Olga's] talk)

Euclidean Wormholes and Cosmologies

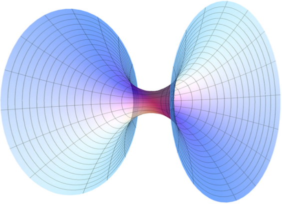
AdS/CFT context: [Maldacena-Maoz (04), PB-Gaddam-Papadoulaki (17) + Kiritsis (19-21), Van Raamsdonk et. al. (20-23) ...]

- Euclidean geometries have interesting connections to Lorentzian geometries via analytic continuation: Slicing them (Z_2 symmetry) we define states/wavefunctions
- Ex: Sphere $\leftrightarrow$ Hartle-Hawking wavefunction (no boundary state) for Cosmology, Euclidean cigar $\leftrightarrow$ (TFD state) for the eternal black hole
- A common misconception: Euclidean Wormholes are NOT related to Black Holes (horizons) via analytic continuation - Instead:



- The most interesting analytic continuation (radial) is related to Cosmologies - these can exist even for negative Λ (Bang/Crunch)
- Along the boundary: leads to traversable wormhole geometries (various known saddles develop pathologies i.e. complex background fields - much harder to obtain negative null energy than negative Euclidean energy)

Euclidean AdS Wormholes (saddles)



- Focus in the case of Euclidean wormhole saddles with Anti-de Sitter asymptotics i.e.

$$ds^2 = d\tau^2 + a^2(\tau) d\Sigma_d^2$$

$$a(\tau = 0) \neq 0\,, \quad a(\tau) \simeq e^{H|\tau|}\,, \quad \tau \to \pm\infty$$

Pertinent Open Questions

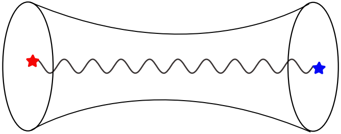
- Microscopic UV complete models of EWs? In AdS/CFT? (we want to understand string theory on target space wormhole backgrounds)
- Worldsheet description of target space wormholes? (putative WZW coset models in [PB - Gaddam - Papadoulaki (2023) + in progress])
- Unfortunately no clear embedding/understanding of such geometries in fully fledged string theory or AdS/CFT so far
- This question is closely related to the factorization problem:

  Entanglement "holds up the throat" of a two sided eternal black hole, but it is not clear what is the analogue for Euclidean wormholes

Symmetries and correlators of local operators

[PB - Kiritsis - Papadoulaki (19), PB - Papadoulaki (23)]

- No obvious entanglement as for Lorentzian wormholes (BH horizons)
- Global symmetries for the boundary theories? $\leftrightarrow$ A common Bulk "Gauss Law constraint" and gauge field
- Symmetries are broken to their diagonal part: $\mathcal{G}_1 \times \mathcal{G}_2 \to \mathcal{G}_{diag.}$ (spontaneous vs. explicit - whether there is a competing factorised saddle or not with the same bcs for bulk fields)

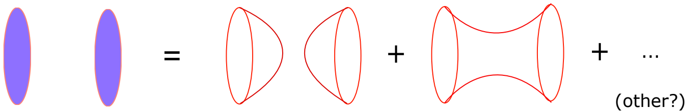


- We find two types of correlation functions, either on a single boundary such as $\langle \mathcal{O}_1 \mathcal{O}_1 \rangle$ or $\langle \mathcal{O}_2 \mathcal{O}_2 \rangle$, or cross-correlators across the two boundaries such as $\langle \mathcal{O}_1 \mathcal{O}_2 \rangle$

- No short distance singularities in the cross-correlators
- One might have expected: different decoupled EQFTs on $\partial \mathcal{M} = \cup_i \partial \mathcal{M}_i$ $\Rightarrow$ Cross correlators are zero or $Z(J_1, J_2) = Z_1(J_1) Z_2(J_2)$
- Wormhole Bulk dictates otherwise $\Rightarrow$ What gives rise to the peculiar properties of the cross-correlators?

The factorisation problem: $Z(J_1, J_2) \neq Z_1(J_1) Z_2(J_2)$

[Maldacena - Maoz (2004) ...]



Possible resolutions in the literature :

- The QGR path integral corresponds to an average: $\langle Z(J_1) Z(J_2) \rangle \Rightarrow$ Several options [...]
- Explicit averaging over ensembles of CFT's - (Unitarity crisis)
- In canonical AdS/CFT there is a single theory with fixed parameters
- Approximate statistical averaging ("ETH" - "Quantum Chaos") $\Rightarrow$ "Statistical wormholes" from complicated/almost random Hamiltonians [...]
- Consistency with $\mathcal{N} = 4$ planar integrability? $\Rightarrow$ Observables/states above the BH threshold [Schlenker - Witten ...]

The "statistical wormholes" need not be saddles of (SU)GRA eoms

No factorisation problem due to interactions?

[PB - Kiritsis - Papadoulaki (19 - 21)], see also related work by [Van Raamsdonk et. al. (20-22)] and [Bachas - Lavdas (18)]

A straightforward resolution for wormhole saddles:

- Interactions between holographic QFT's
- It is actually quite subtle!: "Why to have a disconnected pair of boundaries and not a single one?" ⇒ UV soft - IR strong interactions (reminiscent of confinement...)
- Or: can the exact Schwinger functional acquire an "averaged" form

$$Z_{system}(J_1, J_2) = \sum_S e^{w(S)} Z_S^{(QFT1)}(J_1) Z_S^{(QFT2)}(J_2)$$

  in a single unitary/reflection positive system? (S some "sector")
  [PB - Kiritsis - Papadoulaki (21)]
- Cross correlators ⇒ averages of lower point correlators in individual subsystems - no short distance singularities

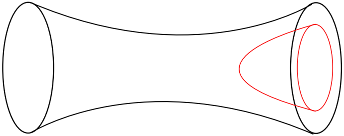
Further (non-local) observables: Wilson Loops

[PB - Kiritsis - Papadoulaki (2019), Refined in: PB - Papadoulaki (2023)]

- Wilson loop observables $W(C) = \mathrm{tr}\left(\mathcal{P} \exp\ i \oint_C A_\mu dx^\mu\right)$ refine the analysis of [Schlenker - Witten (2022)] that studied the compressibility properties of various boundary cycles C in the wormhole bulk

- In holography: Find the string worldsheet ending on the corresponding loop C on a boundary (if it exists) and minimize its area

- Simplest observable: expectation value of a single Wilson loop $\langle W(C) \rangle$



Universal features:

- Large loops on the boundary penetrate further in the bulk and we can probe the IR properties of the boundary dual

- Typically we find an Area law behaviour in the IR

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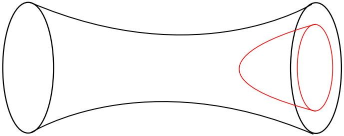

Universal features:

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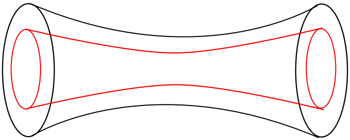
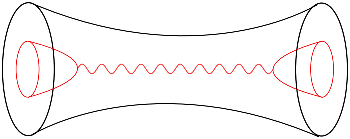
- If the EW geometry contains a non-contractible (thermal) cycle $C_\beta : S^1_\beta$, then there is no bulk surface ending on it, so that $\langle W_P(C_\beta)\rangle = 0$
- Again reminiscent of some kind of confining behaviour (center symmetry) In contrast with the BH cigar for which $\langle W_P(C_\beta)\rangle \neq 0$ (deconfinement)

Wilson Loop correlators (universal results)





- Study loop cross-correlators $\langle W(C_1)W(C_2)\rangle$, the two loops residing on different boundaries
- As we shrink the boundary loops, we find that the leading configuration of lowest action is the one for two disconnected loops
- In the regime of large Wilson loops, the leading contribution originates from a single surface connecting the two loops having a cylinder topology $S^1 \times R$
- Large loops $\Rightarrow$ Strong IR cross-coupling

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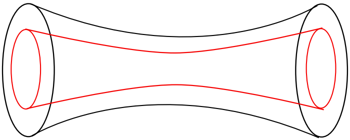

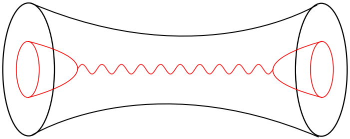

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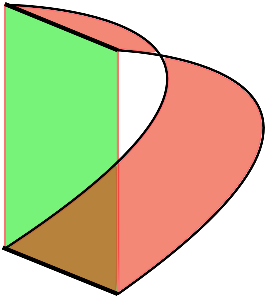
- Large loops $\Rightarrow$ Strong IR cross-coupling
- In the presence of a a non-contractible (thermal) cycle $C_\beta : S^1_\beta$, we find only a connected cylindrical bulk surface ($\langle W_P(C_\beta^{(1)})W_P(C_\beta^{(2)})\rangle \neq 0$)
- Consistent with unbroken diagonal center symmetry ex: $Z_N^{(1)} \times Z_N^{(2)} \to Z_N^{diag.}$ "cross-confining behaviour" - diagonal singlets

Dual QFT models
(reflection positive)

Tripartite BQFT construction

[van Raamsdonk (20) - (22)], [PB - Kiritsis - Papadoulaki (21)]

- Two d-dim (holographic) BQFT's on Σ coupled through a $d+1$-dim intermediate ("messenger") theory on $I \times \Sigma$



- Consider a system for which $c_{d+1} \ll c_d$
- We would like the system to flow to a gapped/confining theory in the IR
- The geometric idea: The dual bulk gravity can localise on $d+1$-dim EOW branes that bend and connect in the IR [van Raamsdonk]

- We focus in the case where the messenger theory is (quasi) topological ($TQFT_{d+1}$) $\Rightarrow$ No contamination from $d+2$ bulk perturbative modes, natural gap in the IR ... [PB - Kiritsis - Papadoulaki]
- Integrate out $TQFT_{d+1}$ $\Rightarrow$ The Schwinger functional does become

$$Z_{system} = \sum_S e^{w(S)} Z_S^{(BQFT_1)}(J_1) Z_S^{(BQFT_2)}(J_2)$$

Solvable microscopic tripartite model $(2d - 1d)$

[PB - Kiritsis - Papadoulaki (21), PB - Papadoulaki (23)]

- Consider a generalised YM in $2d$ (τ, z) with BF action

$$S_{gYM} = \frac{1}{g_{YM}^2}\int_\Sigma \mathrm{tr}\, BF + \frac{\theta}{g_{YM}^2}\int_\Sigma \mathrm{tr}\, B\, d\mu - \frac{1}{2g_{YM}^2}\int_\Sigma \mathrm{tr}\, \Phi(B)\, d\mu$$

  where $F = dA + A\wedge A$

- Couple it with two $1d$ $U(N)$ gauged matrix quantum mechanics theories $M_{1,2}(\tau)$ at the endpoints of an interval I $(z = \pm L)$

$$S_{MQM_{1,2}} = \int d\tau\, \mathrm{tr}\left(\frac{1}{2}(D_\tau M_{1,2})^2 - V(M_{1,2})\right),\ D_\tau M_{1,2} = \partial_\tau M_{1,2} + i[A_\tau^{1,2}, M_{1,2}]$$

  $A_\tau(\tau, z = \pm L) = A_\tau^{1,2}(\tau)$ is the value of the 2d gauge field on the two boundaries

- Solvable system: $2d$ YM - ($\Phi(B) = B^2$) coupled to two Gaussian MQM $(V(M_{1,2}) = \frac{1}{2}M_{1,2}^2)$

"Entangling" the representations (compact $\tau \sim \tau + \beta$)

- Place the system on $I \times S^1$ (cylinder) of length L and circumference β
- The $2d$ YM amplitude on the cylinder is

$$Z_{YM}(U_1, U_2) = \sum_R \chi_R(U_1)\chi_R(U_2^\dagger)e^{-L\frac{g^2_{YM}}{N}C_R^{(2)} + i\theta C_R^{(1)}}$$

and depends on the two asymptotic holonomies $U_{1,2} = \exp \oint d\tau A_\tau^{1,2}$ (zero modes of the gauge field)

- R a $U(N)$ representation, $C_R^{(1,2)}$ its Casimirs and $\chi_R(U)$ are $U(N)$ characters/wavefunctions at the ends of the cylinder
- Integrate out $M_{1,2}$ to obtain the (twisted) MQM partition functions $Z_{1,2}^{MQM}(U_{1,2};\beta) = \int DM_{1,2} \langle U_{1,2}M_{1,2}U_{1,2}^\dagger \,|\, M_{1,2}\rangle_{H.Osc.}$
- Couple the $2d$ YM amplitude $Z_{YM}(U_1, U_2)$ to the two MQM partition functions $Z_{1,2}^{MQM}(U_{1,2};\beta)$ and integrate over the zero modes $U_{1,2}$

"Entangling" the representations (compact $\tau \sim \tau + \beta$)

- The complete partition function on $I \times S^1$ is

$$Z_{system} = \sum_R e^{-L \frac{g_{YM}^2}{N} C_R^{(2)} + i\theta C_R^{(1)}} Z_R^{MQM_1}(\beta) Z_R^{MQM_2}(\beta) \,,$$

$$Z_R^{MQM}(\beta) = \mathrm{tr}_{\mathcal{H}_R} \, e^{-\beta \hat{H}_R^{MQM}} = \int DU \chi_R(U) Z^{MQM}(U;\beta)$$

with β the S^1 size and $\mathcal{H}_R$ the Hilbert space of MQM in a fixed representation R [Kazakov, Klebanov ...]

- The two MQM representations R are correlated/"entangled"

$\sum_R \Rightarrow$ is a form of "averaging", consistent with unitarity (reflection positivity) for a single (tripartite) quantum mechanical system
$\Rightarrow$ What we previously called "the sectors S"

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- No approximation (such as ETH or coarse graining) or averaging over theories involved!

- The allowed representations in the sum are center symmetric (zero n-ality) , so indeed $g_c^{(1)} \times g_c^{(2)} \to g_c^{(diag.)}$ [PB - Papadoulaki (23)]

Higher Dimensional Examples

Sketch of the prescription

- Consider a space with the topology $\mathcal{M} = \Sigma \times I$
- Couple two holographic BQFT's on the interval ends $\Sigma_{1,2}$ using the interval transition amplitude of a (quasi) - topological TQFT in $\Sigma \times I$

2D/3D dimensional example in [PB - Kiritsis - Papadoulaki]

- Two $2D$ BQFT's coupled through a Chern-Simons theory living in $3D$
- Convenient to use radial quantisation ($A_r = 0$) to define the interval transition amplitude [Elitzur-Moore-Schwimmer-Seiberg])
- Couple the CS. transition amplitude to the two $2D$ (gauged) BQFT's and integrate over the gauge field zero modes
- Ex: $\Sigma = T^2$: Replace the wavefunctions $\chi_R(U)$ of the $2D$ YM model with Weyl-Kac characters $\chi_{R,k}(\alpha, \tau)$

- Important to construct a higher dimensional (SUSic) model with control on both sides of the duality ...

$\mathcal{N} = 4$ Wilson loops and type IIB "bubbling wormholes"

Wilson loops in $\mathcal{N}=4$ SYM

- Our $2d/1d$ model is reminiscent of SUSY localization computations of line/defect operators in $\mathcal{N}=4$ SYM [...]
- Our Idea: correlate representations of $(1/2$-BPS) Wilson loops W_R in higher dimensional examples with explicit semiclassical holographic duals. Here: Consider two (non-interacting) copies of $\mathcal{N}=4$ SYM and a correlated observable

$$\sum_R e^{w(R)} \langle W_R \rangle_1 \langle W_R \rangle_2 \qquad W_R = \mathrm{tr}_R\, P \exp\left[i \oint ds (i A_\mu \dot{x}^\mu + \vec{n}\cdot\vec{\Phi}|\dot{x}|)\right]$$

- A single $1/2$-BPS Wilson loop in the representation R is computed via localization resulting in a Hermitean matrix integral [Pestun ...]

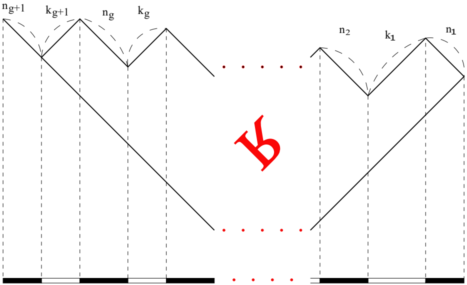
$$\langle W_R \rangle = \langle \mathrm{tr}_R(e^M)\rangle_M = \frac{1}{Z}\int DM e^{-\frac{2N}{\lambda}\,\mathrm{tr}\, M^2} \chi_R\left(e^M\right)$$

- We would like to understand the limit where the operator is "very heavy" and backreacts strongly in the dual geometry

Wilson loops in large ($O(N^2)$) representations

[Gomis, Okuda, Trancanelli ...]

- Representations with $O(N)$ boxes are still light (D-branes etc.)
- We need to consider representations $R : \{R_1, ..R_N\}$ with $O(N^2)$ boxes and the highest weights $R_i \sim O(N)$



- The Young diagram of such reps is described by a collection of rectangular blocks (each with $O(N^2)$ boxes), of size (n_I, k_I) specifying the number of rows/columns
- Once projected onto the real line, a "Maya diagram" is produced consisting of black and white lines

- The lines correspond to the cuts of the matrix model resolvent $\omega_R(z)$, which in turn dictates the form and properties of the dual SUGRA geometry as we shall find out

The type IIB backreacted geometries

- The geometry dual to a backreacted loop in rep R, has an $SO(2,1) \times SO(3) \times SO(5)$ isometry [D'Hoker-Estes- Gutperle, ...]

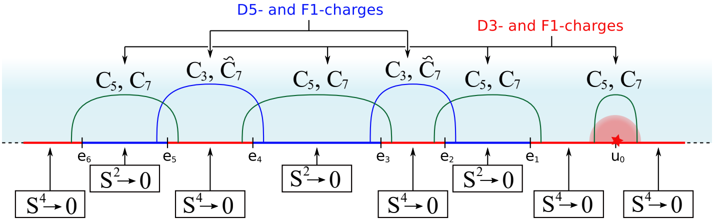
$$ds^2 = f_1^2 ds^2_{AdS_2} + f_2^2 ds^2_{S^2} + f_4^2 ds^2_{S^4} + 4\rho^2 dz d\bar{z}$$

  where $z, \bar{z}$ parametrise a Riemann surface Σ and $f_{1,2,4}(z,\bar{z})$, $\rho(z,\bar{z})$.

- The Wilson loop is on the S^1 boundary of the AdS_2 disk
- The solution also contains a non-trivial dilaton and 3-cycles/5-cycles/7-cycles with RR/RR/NSNS fluxes supporting them $(D5/D3/F1)$

The type IIB backreacted geometries

- The metric, dilaton and fluxes are determined just by two harmonic functions $h_{1,2}(z,\overline{z})$
- $h_2 = 0$ determines the boundary of the Riemann surface Σ (taken to be the upper half-plane)



- $h_1(z,\overline{z}) = \mathcal{A}(z) + \overline{\mathcal{A}}(\overline{z})$ contains all the data of the "bubbling" geometry
- $\mathcal{A}(z)$ cuts $\leftrightarrow$ determine the fluxes and brane charges
  Singularity $\leftrightarrow$ determines the asymptotic $AdS_5 \times S^5$ region

Connecting the MM resolvent with the harmonic functions

- One can show that the matrix model resolvent is related to the two harmonic functions $h_{1,2}$ via ($y(z)$: "spectral - curve")

$$2\omega(z) = V_c'(z) - y(z)\,, \qquad \rho(z) = \frac{1}{2\pi}\Im y(z)\,, \quad z \in \mathcal{C}$$
$$h_1(z,\bar{z}) = \mathcal{A} + \overline{\mathcal{A}}\,, \qquad h_2(z,\bar{z}) = \mathcal{B} + \overline{\mathcal{B}}$$
$$iV_c'(z) = \frac{2i}{\lambda} z = \mathcal{B}(z)\,, \qquad iy(z) = \mathcal{A}(z)$$

- This means that it completely determines the properties of the dual SUGRA geometry

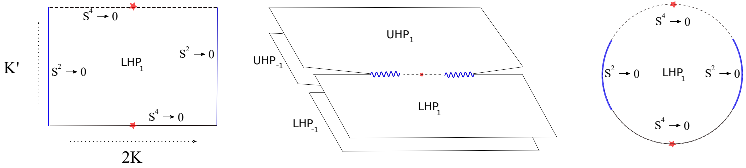
- $h_{1,2}$ need to have common singularities on $\partial\Sigma$. Near such singularities the metric asymptotes to $AdS_5 \times S^5$. ex:

$$h_1 = \frac{2i}{\lambda}\sqrt{z^2 - \lambda} + \text{c.c.}\,, \qquad h_2 = \frac{2i}{\lambda} z + \text{c.c.}$$

- For a single Wilson loop in any rep, there is only a single such singularity. The topology of the boundary is an S^4 and the half-BPS Wilson loop wraps a great $S^1 \subset S^4$

Wormholes $\equiv$ multiple singularities on $\partial\Sigma$

- We found solutions with more than one singularities/asymptotic regions, still preserving the regularity conditions of [D'Hoker-Estes- Gutperle, ...]
- The simplest such Σ corresponds to a disk with two cuts/singularities $\equiv$ a square with two singularities [PB, Ji Hoon Lee, O. Papadoulaki]

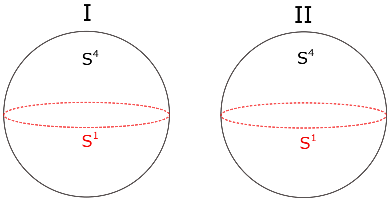


$$h_1(z) = i\frac{2}{\lambda z}\sqrt{(z^2 - e_{min}^2)(z^2 - e_{max}^2)} + cc.\,, \quad h_2(z) = i\frac{2}{\lambda}\left(z - \frac{e_{min}e_{max}}{z}\right) + cc.$$

- We also found more complicated solutions that can be mapped to regular polygons with $2n$ edges and n singularities, as well as solutions when Σ is an annulus

Geometric properties I: AdS_2 factor and "Janus"

- The two boundary wormhole geometry is a form of a double cover of $AdS_5 \times S^5$ (dilaton is still constant)
- There is a caveat: The geometry has an $EAdS_2$ factor with disk topology and its boundary S^1 is shared by all the AdS_5 asymptotic boundaries (Σ singularities) that have the topology of S^4
- This means that the would-be distinct S^4 boundaries are identified on a common S^1, in analogy with other Janus-type of solutions [D'Hoker, Estes, Gutperle, Bachas, Gomis, Assel ...]

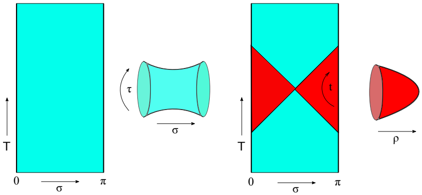


- Still it is possible to connect separate points on the two S^4's by traversing the bulk wormhole, without ever crossing the common S^1

An aside: Two boundary AdS_2 wormhole?

[PB - Papadoulaki (23)]

- What about using global $EAdS_2$ that has two boundaries (cylinder)?



- In this case away from the Σ singularities the geometry is the two boundary $EAdS_2 \times S^4 \times S^2 \times R_2$
(similar to the [Maldacena Milekhin Popov] wormhole geometries)

- At the Σ singularities, the former UV asymptotic S^4's are now replaced by $S^3 \times S^1$

- The two asymptotic S^1's of the cylinder $EAdS_2$ comprise the S^1's on the north and south poles of the S^3.

- Consistent with the fact that one needs to have a pair of Polyakov loops (around the S^1), sitting on the north and south poles of S^3 (Gauss-law)

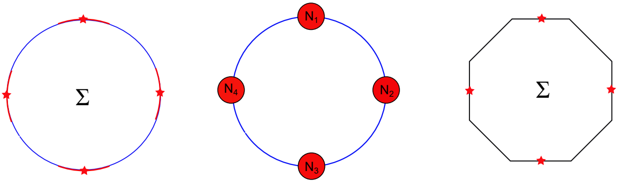
Matrix model dual of Σ wormhole with two S^4 boundaries

- The dual matrix model spectral curve needs two cuts and two singularities
- Use an "analogue of the Dirac-δ" for two $1/2$-BPS loop operators on two copies of $\mathcal{N}=4 \Rightarrow$ "Glue" the two copies of $\mathcal{N}=4$ Wilson loops

$$\langle \det\left(I\otimes I - e^{M_1}\otimes e^{M_2}\right)^{-1}\rangle_{1,2} = \sum_R \langle \chi_R(e^{M_1})\rangle_1 \, \langle \chi_R(e^{M_2})\rangle_2$$

  If the matrices were unitary this would have been a Weyl-invariant delta function

- This can be analysed as a coupled two matrix model or as a model in the space of highest weights R_i of R
- For the multi-boundary wormholes use an $\hat{A}_r$ necklace matrix chain and connect the nodes with determinant operators



Intuitive understanding of the 2MM: Two component gas

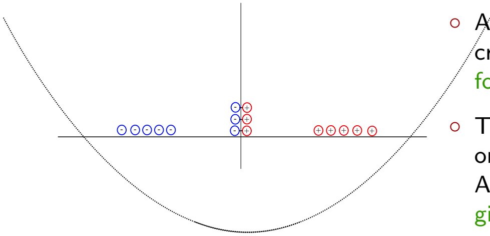
- The 2MM saddle point equations describe two types of particles

$$-\frac{4N_1}{\lambda_1}\mu_i^{(1)} - \sum_{k=1}^{N_2}\frac{2}{\sinh(\mu_i^{(1)}+\mu_k^{(2)})} + \sum_{j\neq i}\frac{2}{\mu_i^{(1)}-\mu_j^{(1)}} = 0\,,$$

$$-\frac{4N_2}{\lambda_2}\mu_k^{(2)} - \sum_{i=1}^{N_1}\frac{2}{\sinh(\mu_i^{(1)}+\mu_k^{(2)})} + \sum_{j\neq k}\frac{2}{\mu_k^{(2)}-\mu_j^{(2)}} = 0$$

  with an $1-1$ and $2-2$ repulsion and $1-2$ attraction to "mirror" points

- There is an overall Gaussian attractive potential ⇒ This leads to a paired $1-2$ condensate at the origin (the additional pole of the planar resolvent)



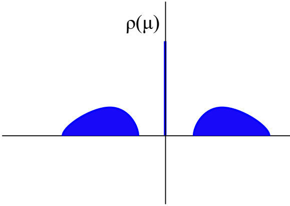
  - After lots of pairs condense, they create a repulsive effective potential for the rest of the eigenvalues
  - The rest of the eigenvalues distribute on two opposite sides of the origin. At large-N they form two cuts, giving rise to the wormhole resolvent

The resolvent at large-N and strong coupling

- At strong 't Hooft coupling the saddle point equations simplify in terms of only rational functions (similar to two coupled $O(2)$ models on a random lattice) [Kostov, Eynard ...]
- In this limit we can obtain an exact solution for the resolvent

$$\omega(z) = \frac{2}{\lambda}\left(z - \frac{ab}{z}\right) - \frac{2}{\lambda z}\sqrt{(z^2 - b^2)(z^2 - a^2)}$$

$$a = \frac{1}{2}(\sqrt{3}-1)\sqrt{\lambda}, \qquad b = \frac{1}{2}(\sqrt{3}+1)\sqrt{\lambda}$$



- The normalisability of the density of eigenvalues ($\int_{supp.} \rho(\mu) = 1$) fixes the end-points a, b in terms of the 't Hooft coupling λ
- The resulting harmonic functions $h_{1,2}$ correspond precisely to the ones we found in the gravitational description

Further properties of wormhole saddle

- One can compare the free energy of the wormhole saddle with two disconnected $AdS_5 \times S^5$ spaces

$$\mathcal{F}_w - 2\mathcal{F}_{AdS} = -\frac{1}{2}\log\lambda$$

- The wormhole has lower free energy. (Indicative for its stability)

- One can also compute the expectation of probe Wilson loops. For example $W_f = \mathrm{tr}\, e^M$

$$\langle W_f \rangle_{AdS} = \int_{-\infty}^{\infty} dz \rho_{AdS}(z) e^z = \frac{2}{\sqrt{\lambda}} I_1(\sqrt{\lambda})$$

$$\langle W_f \rangle_{worm} = \frac{4}{\pi\lambda}\int_a^b \frac{dz}{z}\sqrt{(b^2 - z^2)(z^2 - a^2)}e^z$$

  It grows with a slower rate with λ wtr to the AdS example

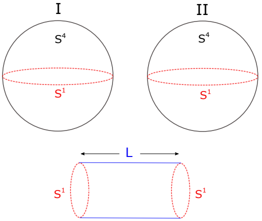
- Interesting to extend this to observables with coordinate dependence, such as correlators of local operators and match with the gravity side

Further comments and generalisations

- Take two copies of $\mathcal{N}=4$ with $\mathcal{G}_1 \times \mathcal{G}_2$ symmetry and consider the general class of correlated observables

$$\sum_R e^{w(R)} \langle \chi_R(e^{M_1}) \rangle_1 \langle \chi_R(e^{M_2}) \rangle_2$$

- $w(R)=0 \Rightarrow$ We "identify" the loops
(ex: $16_1 \times 16_2 \rightarrow 16_{diag}$ - identification of supercharges)



- If we weigh the average with the quadratic Casimir $e^{w(R)} = e^{-LC_R^{(2)}}$ (2d-YM), $\Rightarrow$ the S^1's start to separate (cylinder)
- We can still find the resolvent in this case using the techniques in [Gross-Matytsin, Kazakov ...]
- The density becomes a "time dependent" function $\rho(\mu,\tau)$ obeying the EOMs of collective field theory with appropriate bcs (at $\tau = 0, L$)
- Unfortunately we do not have control over the dual geometry to separate the loops from the bulk side (need a less symmetric ansatze)

Summary and Future

Summary and Future Directions

Summary

- We proposed a general class of microscopic models for Euclidean Wormholes, in terms of BQFTs coupled via a higher dimensional TQFT
- These models are reflection positive and do not require any ad hoc averaging (over couplings/ensembles of CFTs or otherwise)
  - no deviation from the usual holographic prescription and rules

  There is though a resulting sum over representations of the gauge group after we integrate out the "messenger" TQFT
- This makes the resulting field theoretic correlators to be compatible with dual computations on wormhole saddles
- We found that similar models can also arise by considering heavy correlated observables in otherwise decoupled QFTs

  We analysed the case of correlated Wilson loops between copies of $\mathcal{N} = 4$ SYM. They give rise to "bubbling" wormhole geometries in IIB
- In the $1/2$-BPS case we have exact control on both sides of the duality but the boundaries touch on one dimensional $S^1 \subset S^4$'s (similar to Janus)

A Hilbert space interpretation of our constructions

- For Lorentzian wormholes (eternal BH): $\mathcal{H} = \mathcal{H}_{CFT1} \otimes \mathcal{H}_{CFT2}$ and

$$|\Psi\rangle_{TFD} = \sum_n e^{-\frac{\beta}{2} E_n} |E_n\rangle_1 \otimes |E_n\rangle_2$$

- This correlates the energies of the two subsystems
- Our proposed models for Euclidean wormholes: Correlate ("entangle") $U(N)$ representations and not energies as in the TFD
- Realisation I: Presence of gauge constraints (messenger TQFT) - the Hilbert space is reduced into $\mathcal{H} = \sum_R \mathcal{H}_R^1 \otimes \mathcal{H}_R^2$. One could think this in terms of states

$$|\Psi\rangle_{RD} = \sum_R e^{w(R)} |R\rangle_1 \otimes |R\rangle_2$$

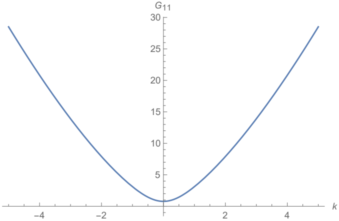
- Realisation II: Consider insertions of "heavy" operators that correlate the copies with a similar representation theoretic "entanglement" (ex: Wilson loops W_R in $\mathcal{N} = 4/IIB$)
- Future Realisation? An effective constraint on the Hilbert space could arise dynamically in the IR ("cross-confinement"/diagonal IR singlets: $U(N) \times U(N) \to U_{diag.}(N)$)

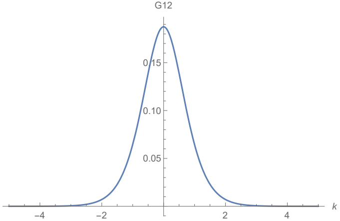
Future Directions

- The MQM non-singlet sectors are also relevant for black hole physics and involve similar sums over representations ($c = 1$ MQM). Connections? [Kazakov et al., PB - Papadoulaki]
- Other top down constructions embeddable in critical string theory ex: Gaiotto Witten systems on an interval [van Raamsdonk, Bachas, ..] Simplify by making the theory on the interval a TQFT "messenger"?
- Less (super)symmetric but still controllable examples of correlated loops or tripartite systems
- Understand better the Lorentzian continuations of our field theoretic setups and their holographic duals (Cosmologies): See [Olga's] talk
- Study (target space) Euclidean wormhole backgrounds in string theory from a worldsheet perspective (WZW cosets?)
- Microscopic "wormhole gas" and replacement of Coleman's $|\alpha\rangle$ states with representations $|R\rangle$ of the dual QFT gauge group?

Thank you!

Scalar Correlators: Universal properties





- Momentum space: $\langle \mathcal{O}_1 \mathcal{O}_1 \rangle$ and $\langle \mathcal{O}_2 \mathcal{O}_2 \rangle$ have a similar behaviour in the UV as in the presence of a single boundary (power law divergence)
- In the IR they saturate to a constant positive value
- The cross correlator $\langle \mathcal{O}_1 \mathcal{O}_2 \rangle$ goes to zero in the UV and has a finite maximum in the IR
- Position space: $(EAdS_2)$ they behave as $\sim 1/\sinh^{2\Delta_+}(\Delta\tau)$ and $\sim 1/\cosh^{2\Delta_+}(\Delta\tau)$ respectively $\Rightarrow$ No short distance singularity for the cross-correlator
- The qualitative behavior of the correlators is similar for several types of solutions $\Rightarrow$ Universality

"Entangling" the representations (infinite τ)

- Define: $J^{\tau}_{MQM_{1,2}} = \delta S_{MQM_{1,2}}/\delta A^{1,2}_{\tau}$ - $U(N)$ MQM charges
- For τ non-compact: $A_\tau = 0 \Rightarrow$ non-perturbative constraint

$$\frac{1}{2g^2_{YM}L} J^{\tau}_{YM} = \frac{1}{2g^2_{YM}L}[W^{-1}, \partial^{\tau} W] = J^{\tau}_{MQM_1} - J^{\tau}_{MQM_2}, \quad W = \mathcal{P}\exp\left(\int_{-L}^{L} dz A_z\right)$$

W a Wilson line extending across the boundaries

- Each MQM Hamiltonian is $(M = U^\dagger \Lambda U,\ J = U^\dagger K U)$

$$\hat{H}^R_{MQM} = \left[-\frac{1}{2}\sum_i \left(\frac{\partial^2}{\partial\lambda_i^2} + V(\lambda_i)\right) + \frac{1}{2}\sum_{i<j}\frac{K^R_{ij}K^R_{ij}}{(\lambda_i - \lambda_j)^2}\right]$$

acting on wave-functions $\Psi_R(\lambda) = \prod_{i<j}(\lambda_i - \lambda_j)\tilde{\Psi}_R(\lambda)$ transforming in the $U(N)$ representation R

- The representations of $MQM_{1,2}$ are "entangled" by the constraint

- The n-point cross-correlator takes the general form

$$\langle O_{i_1}(\tau_{i_1})\dots\tilde{O}_{i_2}(\tau_{i_2})\dots\rangle = \sum_R \langle O_{i_1}(\tau_{i_1})\dots\rangle_1^R \langle\tilde{O}_{i_2}(\tau_{i_2})\dots\rangle_2^R \, e^{-L\frac{g_{YM}^2}{n}C_R^{(2)}+i\theta|R|}$$

where i_1 refers to the first and i_2 to the second MQM subsystem

- This correlator generically only depends separately on the differences $\tau_{i_1}-\tau_{j_1}$ and $\tau_{i_2}-\tau_{j_2}$ and not on time differences that mix the $1,2$ sub-indices, or O_{i_1} with $\tilde{O}_{i_2}$ operators
- No short distance singularities in the cross-correlators!
- The absence of short distance singularities in the cross correlators is a *robust-universal* feature of dual wormhole backgrounds