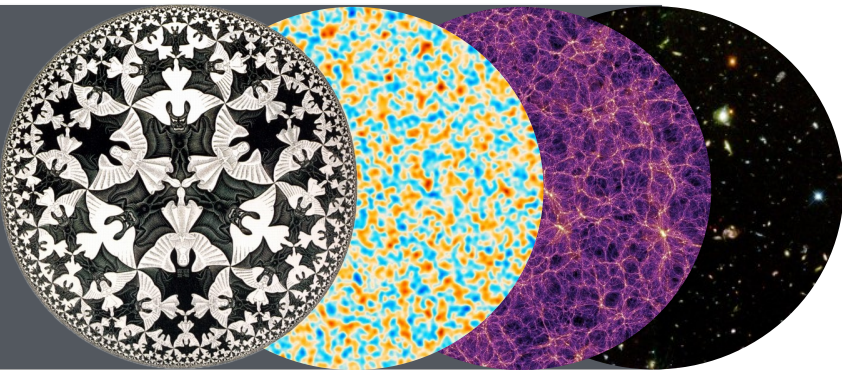


The renormalisation of IR divergences in de Sitter



Paul McFadden (Durham U.)

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Introduction

Holography for cosmology, and in particular for [de Sitter](#), has been explored since the early days of AdS/CFT.

[Strominger, Witten '01; Maldacena '03]

Its status remains ambiguous...

- ▶ Unlike AdS/CFT, we lack a concrete realisation in string theory, and hence examples where both sides of the duality are defined independently.

Nevertheless, supporting evidence exists for a correspondence:

- ▶ Matching of asymptotic symmetries: bulk dS $\leftrightarrow$ boundary Euclidean CFT.
- ▶ Observables defined at future infinity in de Sitter obey the same kinematic constraints as those of a CFT.

While *useful* from a practical standpoint (e.g., for bootstrap approaches), this matching of symmetries is not in itself sufficient to establish a full duality.

Introduction

A similar situation arose in the early days of AdS/CFT where, beyond the matching of symmetries, early support for a full *duality* was provided by holographic renormalisation.

[de Haro, Solodukhin, Skenderis '00]

Gravity in AdS has *IR divergences* due to the infinite volume of spacetime. Via the UV/IR correspondence, these map to *UV divergences* in the dual CFT.

- ▶ In any local dual QFT, these UV divergences are *local*.
Duality requires the same to be true for the bulk IR divergences.
This is a highly nontrivial property satisfied by gravity in AdS.
- ▶ *Renormalisability* of the holographic dictionary requires the existence of appropriate local counterterms on each side of the correspondence.

Introduction

In this talk, we ask:

- ▶ What is the corresponding story in *de Sitter spacetime*?

Cosmological correlators in de Sitter often have IR divergences. At tree level, these arise from the time integrals associated with bulk vertices. We ask:

- ❶ Is the structure of these IR divergences *local*?

This is a *necessary* condition for the existence of a holographic dual description in terms of a local dual CFT.

- ❷ How can IR divergences in de Sitter correlators be renormalised?

We show this is possible by adding *boundary counterterms* defined at future infinity to the cosmological (Schwinger-Keldysh) path integral.

Introduction

This renormalisation procedure, while analogous to holographic renormalisation in AdS, operates **purely in de Sitter** and makes **no use** of holography.

In fact, as we'll discuss, the story in dS is considerably simpler than in AdS.

The renormalisation of IR divergences has practical consequences for the cosmological bootstrap programme:

[Arkani-Hamed, Baumann, Lee, Pimentel '18]

- ▶ IR divergences mean the renormalised de Sitter correlators obey modified, or *inhomogeneous*, conformal Ward identities.
- ▶ A bootstrap approach based on solving the ordinary conformal Ward identities might miss this and fail to reconstruct the correct correlator.

Plan of talk

- ➊ Review of de Sitter correlators.
- ➋ Renormalisation of IR divergences via boundary counterterms.
- ➌ Comparison with holographic renormalisation in AdS.
- ➍ Implications for dS holography.

References

This talk is based on work with [Adam Bzowski](#) and [Kostas Skenderis](#):

- Renormalisation of IR divergences & holography in dS [[JHEP 05 \(2024\) 053](#)]
- A handbook of holographic 4-point functions [[JHEP 12 \(2022\) 039](#)]

Complete results for all 48 renormalised correlators of massless and conformal scalars, at up to 4-points, in both AdS and dS.

See also: Bzowski, Handbook of derivative AdS amplitudes [[JHEP 04 \(2024\) 082](#)]

Other recent work on IR divergences in dS: [[Benincasa & Vazão '24](#)], [[Céspedes, Davis, Wang '23](#)], [[Wang, Pimentel & Achúcarro '22](#)], [[Gorbenko & Senatore '19](#)]

Set-up

- We consider correlators of 'light' scalar fields $\varphi^i(\tau, \mathbf{x})$ with masses

$$m_i^2 = -\Delta_i(\Delta_i - d)L_{dS}^{-2}, \quad d/2 < \Delta_i < d,$$

living on fixed $(d+1)$ -dimensional de Sitter spacetime

$$ds^2 = \frac{L_{dS}^2}{\tau^2}[-d\tau^2 + d\mathbf{x}^2], \quad -\infty < \tau < 0.$$

All interaction vertices will be pure polynomial.

For derivative interactions, see [Bzowski '24]

Set-up

- In the real world, the end of inflation provides a late-time IR cut-off.

Here, a more convenient approach is **dimensional regularisation**:

$$d \rightarrow d + 2\epsilon, \quad \Delta_i \rightarrow \Delta_i + \epsilon.$$

This preserves the de Sitter isometries, and ensures that at late times $\tau \rightarrow 0^-$ the asymptotic fall-off is always power law:

$$\varphi^i(\tau, \mathbf{x}) = (-\tau)^{d-\Delta_i} \varphi_{(0)}^i(\mathbf{x}) + \dots + (-\tau)^{\Delta_i} \varphi_{(\Delta_i)}^i(\mathbf{x}) + \dots$$

- We restrict all calculations to **tree level**.

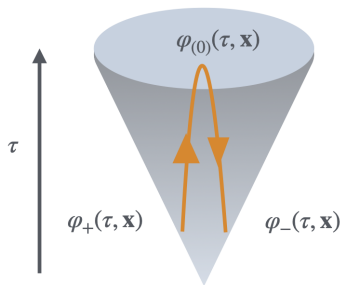
(Loops generate bulk UV divergences which are not cleanly separated from IR divergences in dimensional regularisation.)

In AdS: [Bañados, Bianchi, Muñoz, Skenderis '22]

Correlators in dS

The most direct way to compute dS correlators is via the [Schwinger-Keldysh](#) (or in-in) formalism. This features a closed-time path integral:

$$\begin{aligned} \langle \varphi_{(0)}(\mathbf{x}_1) \dots \varphi_{(0)}(\mathbf{x}_n) \rangle_{dS} &= \int \mathcal{D}\varphi_+(\tau, \mathbf{x}) \mathcal{D}\varphi_-(\tau, \mathbf{x}) \left(\prod_{i=1}^n \varphi_{(0)}(\mathbf{x}_i) \right) \\ &\quad \times \exp \left(iS_+[\varphi_+] - iS_-[\varphi_-] \right) \end{aligned}$$



There are [two](#) distinct fields:

$\varphi_+(\tau, \mathbf{x})$ lives on the 'forwards' and $\varphi_-(\tau, \mathbf{x})$ on the 'backwards' part of the contour.

Differ by $i\epsilon$ prescription: $\tau \rightarrow -\infty(1 \mp i\epsilon)$.

The path integral is restricted to configurations where their late-time values agree:

$$\lim_{\tau \rightarrow 0} (-\tau)^{\Delta-d} \varphi_{\pm}(\tau, \mathbf{x}) = \varphi_{(0)}(\mathbf{x}).$$

Correlators in dS

A **diagrammatic formalism** is set up by introducing **sources** $J_{\pm}$ for the fields $\varphi_{\pm}$:

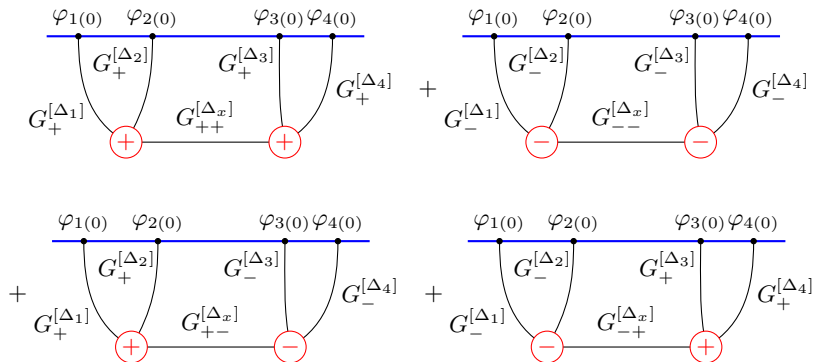
$$Z[J_+, J_-] = \int_{\varphi_+(0, \mathbf{x}) = \varphi_-(0, \mathbf{x}) \sim \varphi_{(0)}(\mathbf{x})} \mathcal{D}\varphi_+ \mathcal{D}\varphi_- \exp \left[iS_+[\varphi_+] - iS_-[\varphi_-] + i \int d^{d+1}x \sqrt{-g} (J_+ \varphi_+ - J_- \varphi_-) \right]$$

- ▶ Vertices come in **two types** $\pm$, according to whether located on the forwards or backwards contour.
- ▶ Bulk-to-bulk propagators then come in **four** types according to the nature of the two vertices they connect: G_{++} , G_{+-} , G_{-+} and G_{--} .
- ▶ Bulk-to-boundary propagators come in **two** types, $G_{\pm}$.
- ▶ Correlators are constructed by summing over diagrams containing all possible assignment of vertex types.

Review: [Chen, Wang, Xianyu 1703.10166]

Correlators in dS

e.g., for a 4-point exchange diagram

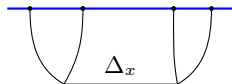


IR divergences

Massless ($\Delta = 3$) and **conformal** ($\Delta = 2$) scalars provide many useful examples where all time integrals can be explicitly evaluated in terms of dilogs, *etc.*

IR divergences are **widespread**: in dim reg, the degrees of divergence ϵ^{-n} are

External Δ_i	3-points
[222]	0
[322]	1
[332]	0
[333]	1



	4-points		
External Δ_i	Contact	$\Delta_x = 2$	$\Delta_x = 3$
[22; 22]	0	0	2
[32; 22]	0	1	1
[33; 22]	1	1	2
[32; 32]	1	2	1
[33; 32]	0	1	1
[33; 33]	1	1	2

Renormalisation

To renormalise the Schwinger-Keldysh path integral

$$Z[J_+, J_-] = \int_{\varphi_+(0, \mathbf{x}) = \varphi_-(0, \mathbf{x}) \sim \varphi_{(0)}(\mathbf{x})} \mathcal{D}\varphi_+ \mathcal{D}\varphi_- \exp \left[iS_+[\varphi_+] - iS_-[\varphi_-] + i \int d^{d+1}x \sqrt{-g} (J_+ \varphi_+ - J_- \varphi_-) \right]$$

we find that a **boundary counterterm** of **just one type** is required:

$$iS_{\text{ct}} = i \int_{\tau=0} d^d \mathbf{x} (J_+ - J_-) f(\varphi_{(0)}; \mu, \mathbf{a}; \epsilon),$$

Here, f is a **local function** of the late-time field $\varphi_{(0)}(\mathbf{x})$ and its derivatives, as well as the RG scale μ and constants $\mathbf{a}$ parametrising the scheme-dependence.

If multiple bulk fields are present, we add such a counterterm for every field, *i.e.*, $(J_+^k - J_-^k) f^k(\{\varphi_0^l\}; \mu, \mathbf{a}_k; \epsilon)$, where the f^k depend on all late-time fields.

Renormalisation

The local functions f^k tells us **how the late-time fields are renormalised**:

At late times, the Schwinger-Keldysh source terms reduce to

$$\int d^{d+1}x \sqrt{-g} (J_+^k \varphi_+^k - J_-^k \varphi_-^k) \rightarrow \int_{\tau=0} d^d \mathbf{x} (J_+^k - J_-^k) \varphi_{(0)}^k$$

and combine with the counterterm to become

$$\int_{\tau=0} d^d \mathbf{x} (J_+^k - J_-^k) \varphi_{R(0)}^k$$

where $\varphi_{R(0)}^k$ is the **renormalised** late-time bulk field given by

$$\varphi_{R(0)}^k = \varphi_{(0)}^k + f^k(\{\varphi_{(0)}^l\}; \mu, \mathfrak{a}_k; \epsilon).$$

Renormalisation

The **form** of the function f^k is constrained by dimensional analysis:

f^k must have dimension $d - \Delta_k$ to match that of the bulk field $\varphi_{(0)}^k$.

Such counterterms exist **precisely** in the cases for which IR divergences arise.

In practice, this means the f^k are polynomials in the late-time fields $\{\varphi_{(0)}^l\}$ and their derivatives. For the n -pt function, we need to know this to degree $(n - 1)$.

For example, in $d = 3$,

- ▶ a massless scalar ($\Delta = 3$) corresponds to a field $\varphi_{(0)}^{[0]}$ of dimension 0.
- ▶ a conformal scalar ($\Delta = 2$) corresponds to a field $\varphi_{(0)}^{[1]}$ of dimension 1.

Renormalisation

Up to 4 points, the available counterterms are:

$$f^{[0]} = a_1^{[0]}(\varphi_{(0)}^{[0]})^2 + a_2^{[0]}(\varphi_{(0)}^{[0]})^3, \quad f^{[1]} = a_1^{[1]}\varphi_{(0)}^{[0]}\varphi_{(0)}^{[1]} + a_2^{[1]}(\varphi_{(0)}^{[0]})^2\varphi_{(0)}^{[1]}$$

The constants $a_j^{[i]}$ are adjusted so that all infinities are cancelled, after which we remove the regulator $\epsilon \rightarrow 0$ to obtain the renormalised correlators.

We showed these counterterms are sufficient to renormalise all 24 IR divergent tree-level correlators of massless and conformal scalars at 3- and 4-points.

Example

Consider, e.g., a massless scalar with $\lambda\varphi^3$ interaction: the renormalised field is

$$\varphi_{(0)}^R = \varphi_{(0)} + \lambda a_2(\mu, \epsilon) : \varphi_{(0)}^2 : + \lambda^2 a_3(\mu, \epsilon) : \varphi_{(0)}^3 :$$

To $O(\lambda^2)$, the renormalised **3-pt function** is

$$\begin{aligned} & \langle \varphi_{(0)}^R \varphi_{(0)}^R \varphi_{(0)}^R \rangle_{ren} \\ &= \lim_{\epsilon \rightarrow 0} \left[\langle \varphi_{(0)} \varphi_{(0)} \varphi_{(0)} \rangle_{reg} + \lambda a_2(\mu, \epsilon) \left(\langle : \varphi_{(0)}^2 : \varphi_{(0)} \varphi_{(0)} \rangle_{reg} + \text{perms} \right) \right] \\ &= \lim_{\epsilon \rightarrow 0} \left[\langle \varphi_{(0)} \varphi_{(0)} \varphi_{(0)} \rangle_{reg} + \lambda a_2(\mu, \epsilon) \left(\langle \varphi_{(0)} \varphi_{(0)} \rangle_{reg} \langle \varphi_{(0)} \varphi_{(0)} \rangle_{reg} + \text{perms} \right) \right] \end{aligned}$$

The momentum-dependence of the ϵ^{-1} divergence in $\langle \varphi_{(0)} \varphi_{(0)} \varphi_{(0)} \rangle_{reg} \propto \lambda$ exactly matches that of the counterterm contribution fixing $a_2(\mu, \epsilon)$ to $O(\epsilon^0)$.

We see this IR divergence is **local**: its structure corresponds to a 4-pt function with two coincident insertions.

Example

At 4 points, inserting the renormalised late time field,

$$\varphi_{(0)}^R = \varphi_{(0)} + \lambda a_2(\mu, \epsilon) : \varphi_{(0)}^2 : + \lambda^2 a_3(\mu, \epsilon) : \varphi_{(0)}^3 :$$

we find

$$\begin{aligned} & \langle \varphi_{(0)}^R \varphi_{(0)}^R \varphi_{(0)}^R \varphi_{(0)}^R \rangle_{ren} \\ &= \lim_{\epsilon \rightarrow 0} \left[\langle \varphi_{(0)} \varphi_{(0)} \varphi_{(0)} \varphi_{(0)} \rangle_{reg} + \lambda a_2(\mu, \epsilon) \left(\langle : \varphi_{(0)}^2 : \varphi_{(0)} \varphi_{(0)} \varphi_{(0)} \rangle_{reg} + \text{perms} \right) \right. \\ & \quad + \lambda^2 a_2^2(\mu, \epsilon) \left(\langle : \varphi_{(0)}^2 : : \varphi_{(0)}^2 : \varphi_{(0)} \varphi_{(0)} \rangle_{reg} + \text{perms} \right) \\ & \quad \left. + \lambda^2 a_3(\mu, \epsilon) \left(\langle : \varphi_{(0)}^3 : \varphi_{(0)} \varphi_{(0)} \varphi_{(0)} \rangle_{reg} + \text{perms} \right) \right] + O(\lambda^3). \end{aligned}$$

As before, the correlators on the RHS now **factorise** into lower-point functions...

Example

Writing $\varphi_{(0)}(x_i) = \varphi_i$ for short, we find

$$\begin{aligned} & \langle \varphi_1^R \varphi_2^R \varphi_3^R \varphi_4^R \rangle_{ren} \\ &= \lim_{\epsilon \rightarrow 0} \left[\langle \varphi_1 \varphi_2 \varphi_3 \varphi_4 \rangle_{reg} + \lambda a_2(\mu, \epsilon) \left(\langle \varphi_1 \varphi_2 \rangle_{reg} \langle \varphi_1 \varphi_3 \varphi_4 \rangle_{reg} + \text{perms} \right) \right. \\ & \quad + \lambda^2 a_2^2(\mu, \epsilon) \left(\langle \varphi_1 \varphi_2 \rangle_{reg} \langle \varphi_1 \varphi_2 \rangle_{reg} \langle \varphi_3 \varphi_4 \rangle_{reg} + \text{perms} \right) \\ & \quad \left. + \lambda^2 a_3(\mu, \epsilon) \left(\langle \varphi_1 \varphi_2 \rangle_{reg} \langle \varphi_1 \varphi_3 \rangle_{reg} \langle \varphi_1 \varphi_4 \rangle_{reg} + \text{perms} \right) \right] + O(\lambda^3) \end{aligned}$$

$\langle \varphi_1 \varphi_2 \varphi_3 \varphi_4 \rangle_{reg}$ has an ϵ^{-2} pole which cancels against the ϵ^{-2} contribution from $\lambda^2 a_2^2(\mu, \epsilon)$. The remaining ϵ^{-1} divergence fixes $a_3(\mu, \epsilon)$ to $O(\epsilon^0)$.

Once again, the IR divergences are **local**: their structure derives from coincident operator insertions, allowing cancellation by a local boundary counterterm.

Renormalisation in AdS

Let's now contrast this procedure with **holographic renormalisation** in AdS.

In AdS, for scalars of mass $m_i^2 = \Delta_i(\Delta_i - d)L_{AdS}^{-2}$ such that $d/2 < \Delta_i < d$, we have the near-boundary asymptotics

$$\begin{array}{ccc} \varphi^i(z, \boldsymbol{x}) = z^{d-\Delta_i} \varphi_{(0)}^i(\boldsymbol{x}) + \dots & + & z^{\Delta_i} \varphi_{(\Delta_i)}^i(\boldsymbol{x}) + \dots \\ \sim \text{source} & & \text{operator } \mathcal{O}_i \end{array}$$

Briefly, the renormalisation of tree-level correlators requires **two steps**:

- ➊ We add **local boundary counterterms** to the regulated on-shell bulk action:

$$S_{\text{on-shell}}[\varphi_{(0)}^i; \epsilon] \rightarrow S_{\text{on-shell}}[\varphi_{(0)}^i; \epsilon] + S_{ct}[\varphi_{(0)}^i; \epsilon]$$

From the dual CFT perspective, these counterterms involve only the CFT **sources**. They encode the contributions from **conformal anomalies**.

[de Haro, Solodukhin & Skenderis '00]

Renormalisation in AdS

- ② We **renormalise the sources**: schematically, up to 3-points, we have

$$\varphi_{(0)}^i = \varphi_{(0)}^i[\phi_{(0)}^j, \epsilon] = \phi_{(0)}^i + \frac{1}{\epsilon} \square^{k_1} \phi_{(0)}^{j_1} \square^{k_2} \phi_{(0)}^{j_2} + \dots$$

This changes the relation between the coefficient $\varphi_{(0)}^i$ in the bulk asymptotic fall-off and the source $\phi_{(0)}^i$ in the dual CFT.

The dimension of the 2nd term here must equal the 1st, but this is precisely the condition for UV singularities in the 3-pt function.

Occurs, e.g., for 3-pt fn of a massless scalar in AdS_4 , see [Bzowski, PM, Skenderis '15]

The renormalised correlators are obtained by functionally differentiating the renormalised on-shell action with respect to these **renormalised** sources $\phi_{(0)}^i$.

$$S_{\text{ren}}[\phi_{(0)}^i] = \lim_{\epsilon \rightarrow 0} \left[S_{\text{on-shell}}[\varphi_{(0)}^i[\phi_{(0)}^j, \epsilon]; \epsilon] + S_{\text{ct}}[\varphi_{(0)}^i[\phi_{(0)}^j, \epsilon]; \epsilon] \right]$$

Renormalisation in AdS

From the dual CFT perspective, we have the coupling

$$S_{\text{CFT}}[\varphi_{(0)}^i, \mathcal{O}_i] = \int d^d \mathbf{x} \varphi_{(0)}^i \mathcal{O}_i.$$

The renormalisation of sources amounts to adding CFT counterterms

$$S_{\text{ct}}[\varphi_{(0)}^i, \mathcal{O}_i; \epsilon] = \frac{1}{\epsilon} \int d^d \mathbf{x} \square^{k_1} \varphi_{(0)}^{j_1} \square^{k_2} \varphi_{(0)}^{j_2} \mathcal{O}_i$$

that renormalise this coupling:

$$\varphi_{(0)}^i = \varphi_{(0)}^i[\phi_{(0)}^j, \epsilon] = \phi_{(0)}^i + \frac{1}{\epsilon} \square^{k_1} \phi_{(0)}^{j_1} \square^{k_2} \phi_{(0)}^{j_2} + \dots$$

While all β -functions vanish at a critical point, their derivatives wrt sources in general do not, and these counterterms encode this data.

[Bzowski, PM, Skenderis '15]

Renormalisation in AdS vs dS

Comparing the asymptotics,

$$\text{AdS: } \varphi^i(z, \mathbf{x}) = z^{d-\Delta_i} \varphi_{(0)}^i(\mathbf{x}) + \dots + z^{\Delta_i} \varphi_{(\Delta_i)}^i(\mathbf{x}) + \dots$$

$\sim$ source operator $\mathcal{O}_i$

$$\text{dS: } \varphi^i(\tau, \mathbf{x}) = (-\tau)^{d-\Delta_i} \varphi_{(0)}^i(\mathbf{x}) + \dots + (-\tau)^{\Delta_i} \varphi_{(\Delta_i)}^i(\mathbf{x}) + \dots$$

late-time field sources $(J_+^i - J_-^i)$

- ▶ The renormalisation of *sources* in AdS is analogous to the renormalisation of *late-time fields* in dS.
- ▶ However, in AdS we also have counterterms $S_{ct}[\{\varphi_{(0)}^i\}; \epsilon]$ involving only CFT sources and generating anomalies. These have *no counterpart* in dS.

Renormalisation in AdS vs dS

The reason why can be seen from the Schwinger-Keldysh path integral:

$$Z[J_+, J_-] = \int_{\varphi_+(0, \mathbf{x}) = \varphi_-(0, \mathbf{x}) \sim \varphi_{(0)}(\mathbf{x})} \mathcal{D}\varphi_+ \mathcal{D}\varphi_- \exp \left[iS_+[\varphi_+] - iS_-[\varphi_-] + i \int d^{d+1}x \sqrt{-g} (J_+ \varphi_+ - J_- \varphi_-) \right]$$

Adding boundary counterterms sends

$$S_+[\varphi_+] \rightarrow S_+[\varphi_+] + S_{ct}[\varphi_{(0)}, J_+], \quad S_-[\varphi_-] \rightarrow S_-[\varphi_-] + S_{ct}[\varphi_{(0)}, J_-].$$

Counterterms of the anomaly type $S_{ct}[\varphi_{(0)}]$ are **independent of the sources $J_{\pm}$** .

They automatically cancel between the forwards and backwards part of the contour, since $Z[J_+, J_-]$ depends only on the **difference** $S_+[\varphi_+] - S_-[\varphi_-]$.

Thus, we can have anomalies in AdS, but not in dS: e.g., the 2-point function in AdS can have logs, but that in dS is always a pure power law.

See also: [Raju et al, '23]

Lessons for dS/CFT

The existence of a holographic duality between dS and a local CFT requires:

- The IR divergences in dS are local (as we showed), but moreover, their **structure** matches that of local UV divergences in the dual CFT.

This last part is **nontrivial**: for specific dimensions $\{\Delta_i, d\}$, CFT correlators have UV divergences associated with anomalies, for which there are **no** corresponding IR divergences in dS.

- Early works explored holographic dictionaries based on **analytic continuation** from AdS to dS. [Maldacena '02], [PM & Skenderis '09]

In these set-ups, anomalies in AdS correlators are automatically projected out in the continuation to dS as a result of their momentum dependence.

Lessons for dS/CFT

More recent works have explored analytic continuations to AdS effective theories featuring fields of the *shadow* dimensions $\bar{\Delta}_i = d - \Delta_i$.

[di Pietro, Gorbenko & Komatsu '21], [Sleight & Taronna '20]

Here, a problem arises: the holographic dictionary is not renormalisable.

Cases arise where the dS correlators are IR divergent and require renormalisation, while the corresponding shadow CFT correlators are finite.

e.g., the dS 3pt fn of two conformal and one massless scalar is *renormalised*

$$\text{ds}_{[322]}^{\text{ren}} = \frac{1}{4q_1^3 q_2 q_3} \left\{ -q_1 + (q_2 + q_3) \left[\log \left(\frac{q_T}{\mu} \right) + \mathfrak{a}_{[322]} - 1 \right] \right\}$$

and does not match the shadow CFT correlator which is *finite*:

$$\langle \mathcal{O}_0 \mathcal{O}_1 \mathcal{O}_1 \rangle_{CFT} = c_{[011]} \frac{(q_2 + q_3)}{q_1^3 q_2 q_3}$$

$\Rightarrow$ only the the scheme-dependent terms in the dS correlator are reproduced.

Conclusions

IR divergences in de Sitter correlators can be removed by adding local counterterms at future infinity to the Schwinger-Keldysh path integral.

- ▶ Only a single type of counterterm is required, which generates a renormalisation of the late-time fields.
- ▶ This is analogous to the renormalisation of sources in AdS/CFT. The anomalies that arise in AdS have no counterpart in dS.

In this sense, renormalisation in dS is simpler than in AdS!

- ▶ We provide explicit results for all 48 renormalised correlators of massless and conformal scalars up to 4-points, in both AdS and dS.
- ▶ Open directions: loops, heavy fields, bootstrapping *renormalised* de Sitter correlators from the associated *inhomogeneous* conformal Ward identities.