

Decoding the string in terms of holographic quantum maps

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It has recently been shown that the Nambu-Goto equation for a string emerges from the junction conditions in three-dimensional gravity. Holographically, gravitational junctions are dual to interfaces in conformal field theory. We demonstrate at the level of linearized gravitational perturbations that each stringy mode of the junction corresponds to a $\mathcal{H}_{\text{in}} \rightarrow \mathcal{H}_{\text{out}}$ quantum map which can be factorized into a scattering matrix involving reflection/transmission and a relative automorphism of the Virasoro algebra, and also a $\mathcal{H}_L \rightarrow \mathcal{H}_R$ map of similar nature. These maps preserve the conformal boundary condition, are independent of the background conformal frame, as in the case of conformal interfaces studied in the literature, and realize a tunable energy transmitter.

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Introduction. The holographic duality states that quantum gravity, particularly a superstring theory, can be reformulated in terms of a nongravitational quantum field theory living at the boundary of spacetime [1–3]. There has been a lot of progress in our understanding of how the bulk spacetime and semiclassical bulk effective theory consisting of Einstein’s gravity coupled to a few fields can be

reconstructed from the boundary quantum field theory in the large N and strong coupling limit [4–7]. The reconstruction of *dynamical* extended objects (branes) of the gravitational theory that are necessary for its non-perturbative completion is therefore a problem of fundamental importance.

Gravitational junctions between three-dimensional asymptotically anti-de Sitter (AdS_3) spacetimes model properties of conformal interfaces in dual two-dimensional conformal field theories (CFTs) [8–11], which appear also in many condensed matter systems, e.g., in quantum wire junctions [12–14], and defects in spin systems [15] and quantum Hall systems [16–18].

The Nambu-Goto equation for the string has recently been shown to emerge from the gravitational junction conditions in three dimensions [19]. In this Letter, we demonstrate at the level of linearized gravitational perturbations how each stringy mode of a gravitational junction gluing two three-dimensional locally AdS_3 spacetimes can

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be decoded as an $\mathcal{H}_{\text{in}} \rightarrow \mathcal{H}_{\text{out}}$ quantum map from the in to out Hilbert spaces of excitations scattered at the dual interface. Equivalently, each stringy mode can be reformulated as a $\mathcal{H}_L \rightarrow \mathcal{H}_R$ quantum map relating the Hilbert spaces of the two dual CFTs that straddle the interface. These maps, which we explicitly construct in the universal sector in the large N limit, preserve the conformal boundary condition, and generalize those realized by conformal interfaces studied in the literature by involving relative automorphisms of the Virasoro algebra. We also prove that the quantum maps realized by the holographic defect operators corresponding to junctions with stringy excitations are independent of the conformal frame (the smoothened dual geometrical background spacetime) at the level of linearized departure from the vacuum (pure AdS_3). This reproduces and generalizes the universality of the reflection/transmission properties of conformal defect operators in CFT that are determined only by their boundary entropies [20].

The conformal interface. The nonvanishing components of the energy-momentum tensor $T_{\mu\nu}$ in any state of a two-dimensional CFT are $T_{\pm\pm} = T_{\mu\nu} n_{\pm}^{\mu} n_{\pm}^{\nu}$, with $n_{\pm}^{\mu}$ being the two future directed null vectors. An interface is a gluing of two CFTs, namely CFT_L and CFT_R on the left and right, respectively. A conformal interface at $x = 0$ is represented by an operator insertion $I_{L,R}$ which satisfies

$$(L_{n,+}^L - L_{n,-}^L)I_{L,R} = I_{L,R}(L_{n,+}^R - L_{n,-}^R), \quad (1)$$

with $L_{n,\pm}^L$ and $L_{n,\pm}^R$ being the Virasoro generators of CFT_L and CFT_R , respectively, as $T_{++} - T_{--}$ generates conformal transformations that leave the line $x = 0$ invariant. Equivalently, folding the right half of spacetime at $x = 0$ and applying reflection on CFT_R , we can represent $I_{L,R}$ as a boundary state $|B\rangle$ in $\text{CFT}_L \otimes \overline{\text{CFT}}_R$ satisfying [11,15,21]

$$(L_{n,+}^L + L_{n,+}^R - L_{n,-}^L - L_{n,-}^R)|B\rangle = 0. \quad (2)$$

In either way, the interface represents a linear map $\mathcal{H}_L \rightarrow \mathcal{H}_R$ from states in CFT_L to states in CFT_R .

The interface can also be viewed as an $\mathcal{H}_{\text{in}} \rightarrow \mathcal{H}_{\text{out}}$ map, e.g., the S matrix relating the incoming right-moving and left-moving energy excitations on the left and right sides of the interface, respectively, and the outgoing left-moving and right-moving energy excitations on the respective left and right sides (see Fig. 1). When both CFTs have the same central charge c , this frequency independent S matrix is

$$\begin{pmatrix} \langle T_{++}^L(\omega) \rangle \\ \langle T_{--}^R(\omega) \rangle \end{pmatrix} = S \begin{pmatrix} \langle T_{--}^L(\omega) \rangle \\ \langle T_{++}^R(\omega) \rangle \end{pmatrix}, \quad S = \begin{pmatrix} \frac{2}{2+\lambda} & \frac{\lambda}{2+\lambda} \\ \frac{\lambda}{2+\lambda} & \frac{2}{2+\lambda} \end{pmatrix}, \quad (3)$$

with $\lambda = \frac{2(c-c_{LR})}{c_{LR}}$, where c_{LR} is the coefficient that appears in the two-point function $\langle T_{++}^L(x) T_{++}^R(y) \rangle_I$ in presence of

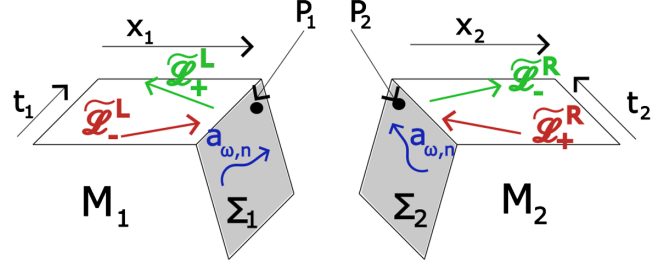


FIG. 1. A two-way gravitational junction formed by gluing two AdS_3 manifolds by identifying points $P_{1,2}$ on the (gray) hypersurfaces. The incoming and outgoing plane wave amplitudes are indicated in the boundary CFT by arrows. The $a_{\omega,n}$ are classical string excitations.

the defect operator [20] and also determines the boundary entropy [22–24] [$T_{\pm\pm}^{L,R}(\omega)$ refers to the Fourier coefficients of $T_{\pm\pm}^{L,R}(x^\pm)$ with $x^\pm = t \pm x$]. Note that S is independent of the frequency in the absence of any dimensionful parameter in the CFTs glued at the interface.

Clearly, S is simply the independent transmission/reflection of the incoming energies from the left and right with equal transmission coefficients

$$\mathbb{T} = \frac{2}{2+\lambda}, \quad (4)$$

and equal reflection coefficients $\mathbb{R} = 1 - \mathbb{T}$, implying energy conservation [$\langle T_{--}^L(\omega) \rangle + \langle T_{++}^R(\omega) \rangle = \langle T_{--}^R(\omega) \rangle + \langle T_{++}^L(\omega) \rangle$], which follows from the conformal boundary condition (1). Remarkably, S is universal for the total energy fluxes, i.e., independent of how the energy fluxes are created [20]. Crucially, an overall (diagonal) conformal transformation on CFT_L and CFT_R does not affect S . The achronal ANEC (averaged null energy condition) implies that $c_{LR} \leq c$ [20], i.e., $\lambda \geq 0$.

Gravitational two-way junction. The conformal interface is holographically dual to a gravitational two-way junction. Consider two three-dimensional locally AdS_3 manifolds $\mathcal{M}_{1,2}$. Each of these is split into two parts $\mathcal{M}_{i\alpha_i}$, $i = 1, 2$, $\alpha_i = L, R$, by codimension-1 hypersurfaces $\Sigma_{1,2}$. A gravitational junction Σ involves the joining of one fragment each of $\mathcal{M}_1$ and $\mathcal{M}_2$. Here we will glue $\mathcal{M}_{1L}$ and $\mathcal{M}_{2R}$. The full spacetime $\tilde{\mathcal{M}}$ is formed by the gluing of $\mathcal{M}_{1L}$ and $\mathcal{M}_{2R}$ at Σ by identifying points on $\Sigma_{1,2}$, which are the images of Σ in $\mathcal{M}_{1,2}$. This identification of points and the embeddings of $\Sigma_{1,2}$ in $\mathcal{M}_{1,2}$ should satisfy the gravitational junction conditions [25].

Let $\mathcal{M}_{1,2}$ have the coordinates $t_{1,2}, x_{1,2}, z_{1,2}$, with $x_{1,2}$ the coordinates transverse to $\Sigma_{1,2}$ and $z_{1,2}$ the radial coordinates. The embeddings of $\Sigma_{1,2}$ are specified by two functions $f_{1,2}$

$$\Sigma_{1,2}: x_{1,2} = f_{1,2}(t_{1,2}, z_{1,2}). \quad (5)$$

The freedom of choice of the coordinates of Σ is fixed by defining the worldsheet coordinates τ and σ at a point P on Σ as

$$\tau(P) = \frac{t_1(P_1) + t_2(P_2)}{2}, \quad \sigma(P) = \frac{z_1(P_1) + z_2(P_2)}{2}, \quad (6)$$

where $P_{1,2}$ are the points on $\Sigma_{1,2}$ that are identified with P (see Fig. 1). The following four variables, which are functions of τ and σ , completely specify the junction

$$\tau_d = \frac{t_2 - t_1}{2}, \quad \sigma_d = \frac{z_2 - z_1}{2}, \quad x_{s,d} = \frac{f_2 \pm f_1}{2}. \quad (7)$$

The junction conditions can be obtained from the action

$$S_{\text{grav}} = \frac{1}{16\pi G_N} \int_{\mathcal{M}} d^3x \sqrt{-g} (R - 2\Lambda) + T_0 \int_{\Sigma} d\tau d\sigma \sqrt{-\gamma} + \text{GHY terms}, \quad (8)$$

where the metric is the *only* degree of freedom. Above GHY are the Gibbons-Hawking-York (GHY) boundary terms and T_0 is the tension. This action is defined assuming that the induced metrics $\gamma_{1,2}$ on $\Sigma_{1,2}$ are identical, which defines the worldsheet metric as

$$\gamma_{\mu\nu}(\tau, \sigma) = \gamma_{1,\mu\nu}(\tau, \sigma) = \gamma_{2,\mu\nu}(\tau, \sigma). \quad (9)$$

Varying (8) away from the junction implies that the manifold is Einstein. At the junction, we have

$$(K_{i,\mu\nu} - K_i \gamma_{i,\mu\nu})|_{\text{disc}} = 8\pi G_N T_0 \gamma_{\mu\nu}, \quad (10)$$

where $K_{i,\mu\nu}$ is the extrinsic curvature of Σ_i in $\mathcal{M}_{i\alpha_i}$, $K_i = K_{i,\mu\nu} \gamma^{\mu\nu}$, and $|_{\text{disc}}$ denotes the discontinuity. The bulk diffeomorphism symmetry implies that the left-hand side of (10) is conserved. We therefore obtain only one independent equation from (10), which together with (9) yields four equations for the four unknown functions (7).

It has been shown in [19], when $\mathcal{M}_1$ and $\mathcal{M}_2$ are copies of an Einstein manifold $\mathcal{M}$, that the general solutions of the junction conditions are in one-to-one correspondence with solutions of the nonlinear Nambu-Goto (NG) equation for a worldsheet in $\mathcal{M}$. Specifically, the hypersurface

$$\Sigma_{NG}: t = \tau, \quad z = \sigma, \quad x = x_s(\tau, \sigma)$$

corresponds to a solution of the nonlinear Nambu-Goto equation in $\mathcal{M}$ when the tension T_0 vanishes, while τ_d , σ_d , and x_d are fixed completely by x_s up to six rigid parameters related to worldsheet and spacetime isometries, which are irrelevant for the present paper.

Here we will focus on locally AdS_3 Bañados metrics [26]

$$ds^2 = \frac{dz^2}{z^2} + 2dt dx (\mathcal{L}_+(x^+) - \mathcal{L}_-(x^-)) - \frac{dt^2}{z^2} (1 - z^2 \mathcal{L}_+(x^+))(1 - z^2 \mathcal{L}_-(x^-)) + \frac{dx^2}{z^2} (1 + z^2 \mathcal{L}_+(x^+))(1 + z^2 \mathcal{L}_-(x^-)), \quad (11)$$

where $x^\pm = t \pm x$, and we have set $\Lambda = -1$. We assume that $\mathcal{M}_{1,2}$ have the above metrics with

$$\mathcal{L}_\pm^{(1)}(x_1^\pm) = \mathcal{L}_{\omega,\pm}^L e^{i\omega x_1^\pm}, \quad \mathcal{L}_\pm^{(2)}(x_2^\pm) = \mathcal{L}_{\omega,\pm}^R e^{i\omega x_2^\pm}. \quad (12)$$

Furthermore, we assume that $\mathcal{L}_{\omega,\pm}^{L,R}$ are small amplitudes of $\mathcal{O}(\epsilon)$. Our analysis will be linear and the above plane waves can be superposed to form wave packets. At $\mathcal{O}(\epsilon^0)$ we have the exact solution

$$\tau_d = 0, \quad \sigma_d = 0, \quad x_s = 0, \quad x_d = -\frac{\lambda\sigma}{\sqrt{4-\lambda^2}}, \quad (13)$$

where $\lambda = 8\pi G_N T_0$ and $0 \leq \lambda \leq 2$. Equations (9) and (10) can be solved perturbatively in ϵ . At $\mathcal{O}(\epsilon)$ we get the following equation:

$$16\sigma\partial_\tau^2 x_s + 4(4-\lambda^2)(2\partial_\sigma x_s - \sigma\partial_\sigma^2 x_s) = -ie^{i\omega\tau}(-4+\lambda^2)\sigma^3\omega\left((\mathcal{L}_{\omega,-}^R - \mathcal{L}_{\omega,+}^L)e^{\frac{i\lambda\omega\sigma}{\sqrt{4-\lambda^2}}} + (\mathcal{L}_{\omega,-}^L - \mathcal{L}_{\omega,+}^R)e^{-\frac{i\lambda\omega\sigma}{\sqrt{4-\lambda^2}}}\right), \quad (14)$$

which is just the linearized Nambu-Goto equation in empty AdS with sources proportional to $(\mathcal{L}_{\omega,-}^L - \mathcal{L}_{\omega,+}^R)$ and $(\mathcal{L}_{\omega,-}^R - \mathcal{L}_{\omega,+}^L)$ when $\lambda \rightarrow 0$. This implies that the correspondence between the Nambu-Goto equation and junction conditions shown in [19] generalizes (with sources) even when the backgrounds on both sides depart from each other. We obtain the following solution to (14):

$$\begin{aligned}
 x_s = e^{i\omega\tau} & \left(\frac{-i(4-\lambda^2)^{1/4}(\sin(\frac{2\sigma\omega}{\sqrt{4-\lambda^2}})(A_{\omega,1}\sqrt{4-\lambda^2} + 2A_{\omega,2}\sigma\omega) + \cos(\frac{2\sigma\omega}{\sqrt{4-\lambda^2}})(A_{\omega,2}\sqrt{4-\lambda^2} - 2A_{\omega,1}\sigma\omega))}{2\sqrt{\pi}(\omega)^{3/2}} \right. \\
 & + \frac{(\mathcal{L}_{\omega,-}^R - \mathcal{L}_{\omega,+}^L)e^{\frac{i\lambda\sigma\omega}{\sqrt{4-\lambda^2}}}(2\sqrt{4-\lambda^2}\lambda\sigma\omega - i(\lambda^2-4)(\sigma^2\omega^2+2))}{4(\lambda^2-4)\omega^3} \\
 & \left. + \frac{(\mathcal{L}_{\omega,-}^L - \mathcal{L}_{\omega,+}^R)e^{-\frac{i\lambda\sigma\omega}{\sqrt{4-\lambda^2}}}(-2\sqrt{4-\lambda^2}\lambda\sigma\omega - i(\lambda^2-4)(\sigma^2\omega^2+2))}{4(\lambda^2-4)\omega^3} \right) + \mathcal{O}(\epsilon^2), \quad (15)
 \end{aligned}$$

where the first line is the solution of the source-free (homogeneous) Eq. (14). Imposing ingoing boundary conditions at the Poincaré horizon [27–29] we obtain the coefficients $A_{\omega,1} = A_{\omega,nn} + A_{\omega,n}$ and $A_{\omega,2} = iA_{\omega,nn}$ which are of $\mathcal{O}(\epsilon)$. Here $A_{\omega,nn}$ corresponds to a non-normalizable mode of the homogeneous NG equation, which is the causal response to bulk perturbations that travels from the boundary toward the Poincaré horizon. $A_{\omega,n}$ is an intrinsic normalizable (stringy) mode of the homogeneous NG equation. Both $A_{\omega,nn}$ and $A_{\omega,n}$ are determined by initial and boundary conditions as usual in Lorentzian holographic duality. Explicit solutions for the other variables are in the Supplemental Material [30].

The energy-momentum tensors on both sides of the dual CFT interface can be extracted using holographic renormalization [31,32]. Explicitly,

$$\langle T_{\pm\pm}^L(\omega) \rangle = \frac{c}{12\pi} \mathcal{L}_{\omega,\pm}^L, \quad \langle T_{\pm\pm}^R(\omega) \rangle = \frac{c}{12\pi} \mathcal{L}_{\omega,\pm}^R. \quad (16)$$

The Dirichlet boundary conditions corresponding to the interface at $x_1 = x_2 = 0$ impose $\lim_{\sigma \rightarrow 0} x_d = \lim_{\sigma \rightarrow 0} x_s = 0$. From $\lim_{\sigma \rightarrow 0} x_d = 0$ we obtain

$$\mathcal{L}_{\omega,-}^L + \mathcal{L}_{\omega,+}^R = \mathcal{L}_{\omega,-}^R + \mathcal{L}_{\omega,+}^L. \quad (17)$$

This is the conservation of energy at the interface, or equivalently the conformal boundary condition. Equation (17) can be solved by

$$\begin{aligned}
 \mathcal{L}_{\omega,+}^L &= \mathcal{T}_{\omega,L} \mathcal{L}_{\omega,+}^R + (1 - \mathcal{T}_{\omega,R}) \mathcal{L}_{\omega,-}^L, \\
 \mathcal{L}_{\omega,-}^R &= \mathcal{T}_{\omega,R} \mathcal{L}_{\omega,-}^L + (1 - \mathcal{T}_{\omega,L}) \mathcal{L}_{\omega,+}^R, \quad (18)
 \end{aligned}$$

where $\mathcal{T}_{\omega,L(R)}$ and $1 - \mathcal{T}_{\omega,L(R)}$ are arbitrary (frequency-dependent) transmission and reflection coefficients for the left and right movers, respectively [note these are $\mathcal{O}(\epsilon^0)$]. From $\lim_{\sigma \rightarrow 0} x_s = 0$, we obtain that

$$A_{\omega,nn} = i \frac{2\sqrt{\pi}(\mathcal{L}_{\omega,-}^L - \mathcal{T}_{\omega,R} \mathcal{L}_{\omega,-}^R - \mathcal{L}_{\omega,+}^R + \mathcal{T}_{\omega,L} \mathcal{L}_{\omega,+}^L)}{(4 - \lambda^2)^{3/4} \omega^{3/2}}. \quad (19)$$

The string as a holographic quantum map. In [33], it was shown that the transmission/reflection coefficients of the CFT interface can be derived by keeping the incoming energy flux only on one side, considering the causal response on the worldsheet with $A_{\omega,n} = 0$ and using the Dirichlet boundary conditions on the other variables τ_d and σ_d . However, the gravitational problem is fundamentally nonlinear, and therefore assuming independent and equal transmission from the left and right sides affects the higher-order perturbative expansions. Furthermore, as discussed below, setting Dirichlet boundary conditions for τ_d and σ_d is inadequate for proving that the scattering process is independent of the conformal frame, i.e., the smoothened dual background geometry in the absence of the interface (which is assumed to be the empty AdS_3 space dual to the CFT vacuum on both sides).

The key point is that, instead of setting Dirichlet boundary conditions for τ_d and σ_d , we should use relative conformal transformations to obtain continuous coordinates and metric across the dual interface at the boundary in which physical experiments are set up. Let us define

$$\lim_{\sigma \rightarrow 0} \tau_d(\tau, \sigma) = \mathbb{t}_{\omega,d} e^{i\omega\tau} + \mathcal{O}(\epsilon^2) \quad (20)$$

so that at the boundary $t_{1,2}(\tau) = \tau \mp \mathbb{t}_{\omega,d} e^{i\omega\tau} + \mathcal{O}(\epsilon^2)$. In agreement with [19], we find that the relative time reparametrization $\mathbb{t}_d(\tau)$ encodes the normalizable stringy mode as both $\mathbb{t}_{\omega,d}$ and $\lim_{\sigma \rightarrow 0} \frac{\sigma_d(\tau, \sigma)}{\sigma}$ are proportional to

$$\begin{aligned}
 & -\mathcal{L}_{\omega,-}^L(-2 + \mathcal{T}_{\omega,R}(2 + \lambda)) + \mathcal{L}_{\omega,+}^R(-2 + \mathcal{T}_{\omega,L}(2 + \lambda)) \\
 & + 2a_{\omega,n},
 \end{aligned}$$

where $a_{\omega,n} = \frac{iA_{\omega,n}\lambda(4-\lambda^2)^{3/4}\omega^{3/2}}{4\sqrt{\pi}}$.

The discontinuity in the time coordinates (t_1 and t_2) at the interface located at $x_1 = x_2 = 0$ can be undone using separate conformal transformations on the two sides which involve coordinate transformations,

$$\begin{aligned}\tilde{t}_{1,2} &= \frac{1}{2}(\mathbb{h}_{1,2}^{-1}(t_{1,2} + x_{1,2}) + \mathbb{h}_{1,2}^{-1}(t_{1,2} - x_{1,2})), \\ \tilde{x}_{1,2} &= \frac{1}{2}(\mathbb{h}_{1,2}^{-1}(t_{1,2} + x_{1,2}) - \mathbb{h}_{1,2}^{-1}(t_{1,2} - x_{1,2})),\end{aligned}\quad (21)$$

and the associated Weyl transformations that bring the metric on both sides back to the Minkowski form. The most general choices of $\mathbb{h}_{1,2}$ at the linear order are

$$\mathbb{h}_2(\tau) = \tau + \kappa_\omega \mathbb{t}_{\omega,d} e^{i\omega\tau}, \quad \mathbb{h}_1(\tau) = \tau + (\kappa_\omega - 2) \mathbb{t}_{\omega,d} e^{i\omega\tau}, \quad (22)$$

where κ_ω is an arbitrary $\mathcal{O}(\epsilon^0)$ (frequency-dependent) parameter which acts as an overall (diagonal) conformal transformation that changes the conformal frame (see below). It is easy to see using (22) that (21) preserves the interface at $\tilde{x}_{1,2} = 0$ where $\tilde{t}_2 = \tilde{t}_1$. We therefore obtain continuous coordinates and metric across the interface. These conformal transformations can be uplifted to bulk diffeomorphisms $[X^\mu(\tau, \sigma) \rightarrow \tilde{X}^\mu(\tau, \sigma)]$ with X^μ denoting bulk coordinates on both sides [29,32,34], and both the induced metric and the extrinsic curvatures of $\Sigma_{1,2}$ remain invariant under these bulk diffeomorphisms, producing an equivalent solution of the junction conditions.

Under the conformal transformations, the energy-momentum tensors on both sides of the interface transform as

$$\begin{aligned}\tilde{T}_{\pm\pm}^{L,R}(\tilde{x}^\pm) &= \mathbb{h}'_{1,2}(\tilde{x}^\pm)^2 T_{\pm\pm}^{L,R}(\mathbb{h}_{1,2}(\tilde{x}_2^\pm)) \\ &\quad - \frac{c}{24\pi} \text{Sch}(\mathbb{h}_{1,2}(\tilde{x}^\pm), \tilde{x}^\pm) \\ &= \frac{c}{12\pi} e^{i\omega\tilde{x}_2^\pm} \tilde{\mathcal{L}}_{\omega,\pm}^{L,R} + \mathcal{O}(\epsilon^2),\end{aligned}\quad (23)$$

where

$$\begin{aligned}\tilde{\mathcal{L}}_{\omega,+}^L &= \frac{\mathcal{L}_{\omega,+}^R}{4} (4 + 2(-1 + \mathcal{T}_{\omega,L})\kappa_\omega + \mathcal{T}_{\omega,L}(-2 + \kappa_\omega)\lambda) \\ &\quad + \frac{\mathcal{L}_{\omega,-}^L}{4} (2\kappa_\omega + \mathcal{T}_{\omega,R}(2\lambda - \kappa_\omega(2 + \lambda))) + \frac{a_{\omega,n}(\kappa_\omega - 2)}{2},\end{aligned}\quad (24)$$

$$\begin{aligned}\tilde{\mathcal{L}}_{\omega,-}^L &= \frac{\mathcal{L}_{\omega,+}^R}{4} (\kappa_\omega - 2)(-2 + \mathcal{T}_{\omega,L}(2 + \lambda)) \\ &\quad + \frac{\mathcal{L}_{\omega,-}^L}{4} (2\kappa_\omega - \mathcal{T}_{\omega,R}(-2 + \kappa_\omega)(2 + \lambda)) + \frac{a_{\omega,n}(\kappa_\omega - 2)}{2},\end{aligned}\quad (25)$$

$$\begin{aligned}\tilde{\mathcal{L}}_{\omega,+}^R &= \frac{\mathcal{L}_{\omega,+}^R}{4} (4 + \kappa_\omega(-2 + \mathcal{T}_{\omega,L}(2 + \lambda))) \\ &\quad - \frac{\mathcal{L}_{\omega,-}^L}{4} \kappa_\omega(-2 + \mathcal{T}_{\omega,R}(2 + \lambda)) + \frac{\kappa_\omega a_{\omega,n}}{2},\end{aligned}\quad (26)$$

$$\begin{aligned}\tilde{\mathcal{L}}_{\omega,-}^R &= \frac{\mathcal{L}_{\omega,+}^R}{4} (2(-1 + \mathcal{T}_{\omega,L})(-2 + \kappa_\omega) + \mathcal{T}_{\omega,L}\kappa_\omega\lambda) \\ &\quad + \frac{\mathcal{L}_{\omega,-}^L}{4} (2\kappa_\omega + \mathcal{T}_{\omega,R}(4 - \kappa_\omega(2 + \lambda))) + \frac{\kappa_\omega a_{\omega,n}}{2}.\end{aligned}\quad (27)$$

These transformations are reproduced in the bulk via the holographic renormalization procedure after implementing the bulk diffeomorphisms which uplift the corresponding conformal transformations [34]. The transformed amplitudes also satisfy energy conservation (and thus also conformal boundary condition) as

$$\tilde{\mathcal{L}}_{\omega,-}^L + \tilde{\mathcal{L}}_{\omega,+}^R = \tilde{\mathcal{L}}_{\omega,-}^R + \tilde{\mathcal{L}}_{\omega,+}^L. \quad (28)$$

As is evident from (24), (25), (26), and (27), the four physical energy fluxes $\tilde{\mathcal{L}}_{\omega,\pm}^{L,R}$ depend on the six parameters κ_ω , $a_{\omega,n}$, $\mathcal{L}_{\omega,+}^R$, $\mathcal{L}_{\omega,-}^L$, $\mathcal{T}_{\omega,L}$, and $\mathcal{T}_{\omega,R}$. Remarkably, for arbitrary values of the six parameters, we obtain that the physical excitations satisfy

$$\begin{aligned}\tilde{\mathcal{L}}_{\omega,+}^L &= \frac{2}{2 + \lambda} \left(\tilde{\mathcal{L}}_{\omega,+}^R - \frac{2a_{\omega,n}}{2 - \lambda} \right) + \frac{\lambda}{2 + \lambda} \left(\tilde{\mathcal{L}}_{\omega,-}^L + \frac{2a_{\omega,n}}{2 - \lambda} \right), \\ \tilde{\mathcal{L}}_{\omega,-}^R &= \frac{\lambda}{2 + \lambda} \left(\tilde{\mathcal{L}}_{\omega,+}^R - \frac{2a_{\omega,n}}{2 - \lambda} \right) + \frac{2}{2 + \lambda} \left(\tilde{\mathcal{L}}_{\omega,-}^L + \frac{2a_{\omega,n}}{2 - \lambda} \right).\end{aligned}\quad (29)$$

This implies that the dual interface acts as an $\mathcal{H}_{\text{in}} \rightarrow \mathcal{H}_{\text{out}}$ quantum map in the universal sector given by $\mathcal{S} \circ \mathcal{D}$ where $\mathcal{S}$ is the matrix (3) and $\mathcal{D}$ is a redistribution of energy among the two incoming energy excitations, which can be realized by the linearized conformal transformation given by

$$g_\pm(x^\pm) = x^\pm \mp \frac{e^{i\omega x^\pm}}{\omega^3} \frac{4ia_{\omega,n}}{2 - \lambda} + \mathcal{O}(\epsilon^2). \quad (30)$$

Crucially, the quantum map $\mathcal{S} \circ \mathcal{D}$ is independent of κ_ω which changes all $\tilde{\mathcal{L}}_{\omega,\pm}^{L,R}$ by an additive factor of $\kappa_\omega a_{\omega,n}/2$ and thus the overall conformal frame (i.e., the smoothened background geometry) reproducing and generalizing this feature of static conformal interfaces [20]. The parameter $a_{\omega,n}$ determining the energy redistribution $\mathcal{D}$ can be extracted from $\langle D(\omega) \rangle$, the expectation value of the displacement operator (which is also independent of κ_ω) as shown using the Ward identities (see Supplemental material [30], which includes Ref. [35], for a derivation).

The physical excitations (24), (25), (26), and (27) also satisfy

$$\begin{aligned}\tilde{\mathcal{L}}_{\omega,+}^L &= \frac{2}{2+\lambda}\tilde{\mathcal{L}}_{\omega,+}^R + \frac{\lambda}{2+\lambda}\tilde{\mathcal{L}}_{\omega,-}^L - \frac{2a_{\omega,n}}{2+\lambda}, \\ \tilde{\mathcal{L}}_{\omega,-}^R &= \frac{\lambda}{2+\lambda}\tilde{\mathcal{L}}_{\omega,+}^R + \frac{2}{2+\lambda}\tilde{\mathcal{L}}_{\omega,-}^L + \frac{2a_{\omega,n}}{2+\lambda},\end{aligned}\quad (31)$$

so that the $\mathcal{H}_{\text{in}} \rightarrow \mathcal{H}_{\text{out}}$ map can be rewritten in the form $\tilde{\mathcal{D}} \circ \mathcal{S}$ where $\tilde{\mathcal{D}}$ is a redistribution of energy among the two outgoing energy excitations which can also be realized by a linearized conformal transformation.

We also note from (24), (25), (26), and (27) that

$$\begin{aligned}\tilde{\mathcal{L}}_{\omega,+}^R &= \left(1 + \frac{\lambda}{2}\right)\tilde{\mathcal{L}}_{\omega,+}^L - \frac{\lambda}{2}\tilde{\mathcal{L}}_{\omega,-}^L + a_{\omega,n}, \\ \tilde{\mathcal{L}}_{\omega,-}^R &= \frac{\lambda}{2}\tilde{\mathcal{L}}_{\omega,+}^L + \left(1 - \frac{\lambda}{2}\right)\tilde{\mathcal{L}}_{\omega,-}^L + a_{\omega,n}.\end{aligned}\quad (32)$$

Thus the interface acts as a universal $\mathcal{H}_L \rightarrow \mathcal{H}_R$ map of the form $\mathcal{C} \circ \tilde{\mathcal{S}}$ mapping the energy fluxes on left side to those on the right side with

$$\tilde{\mathcal{S}} = \begin{pmatrix} 1 + \frac{\lambda}{2} & -\frac{\lambda}{2} \\ \frac{\lambda}{2} & 1 - \frac{\lambda}{2} \end{pmatrix} \quad (33)$$

being an invertible matrix and $\mathcal{C}$ is a conformal transformation, determined by $a_{\omega,n}$, on the right side that acts in the same way on the left and right movers.

For any value of λ between 0 and 2, the interface is topological when

$$a_{\omega,n} = \frac{\lambda}{2}(\tilde{\mathcal{L}}_{\omega,-}^L - \tilde{\mathcal{L}}_{\omega,+}^R). \quad (34)$$

In this case, as is evident from (29), $\tilde{\mathcal{L}}_{\omega,\pm}^R = \tilde{\mathcal{L}}_{\omega,\pm}^L$ so that we obtain perfect energy transmission. Similarly, the interface is always factorizing when

$$a_{\omega,n} = \tilde{\mathcal{L}}_{\omega,+}^R - \tilde{\mathcal{L}}_{\omega,-}^L \quad (35)$$

as in this case $\tilde{\mathcal{L}}_{\omega,+}^{R,L} = \tilde{\mathcal{L}}_{\omega,-}^{R,L}$ and we obtain perfect energy reflection. Thus remarkably, the stringy interface is completely tunable. As discussed in the Supplemental Material [30], the ANEC is not violated as the additional conformal transformation factors do not affect the total energy fluxes.

Generalized conformal interfaces. Our calculations up to the linear order in the energy excitations reveal that the holographic interface operator corresponding to a stringy junction mode is a modification of (1) taking the form

$$(L_{n,+}^L - L_{-n,-}^L)I_{L,R} = I_{L,R}(\tilde{L}_{n,+}^R - \tilde{L}_{-n,-}^R), \quad (36)$$

implying (28), where $\tilde{L}_{m,\pm}$ are automorphisms of $L_{\pm,m}$ which preserve the Virasoro algebra. Although we have not computed the gravitational perturbations beyond the

linearized order, our result addresses a key conceptual issue relevant even for the static holographic junction studied in [33] at the linearized level. Nonlinearity of gravity, particularly the junction conditions, implies mixing of $\mathcal{L}_{2\omega,\pm}^{L,R}$ with $\mathcal{L}_{\omega,\pm}^{L,R^2}$ for instance, while the dual interface should realize a linear quantum map. Consider an automorphism of the form $L_n \rightarrow \tilde{L}_n = VL_n V^{-1}$ with

$$V = \exp \left(\sum_m \zeta_m L_m + \sum_{k=2}^{\infty} \sum_{p_1} \cdots \sum_{p_k} \tilde{\zeta}_{p_1, \dots, p_k} L_{p_1} \cdots L_{p_k} \right).$$

When $\tilde{\zeta}_{p_1, \dots, p_k}$ vanish, the above is a (non)unitary conformal transformation which acts linearly on L_n . However, generally the automorphism acts nonlinearly on L_n in spite of being a linear quantum map. So (36) is consistent with the nonlinearity of the gravitational problem.

At the level of linearized conformal transformations, the factors $\mathcal{D}$, $\tilde{\mathcal{D}}$, and $\mathcal{C}$ in the quantum maps discussed above are conformal transformations determined only by the homogeneous Nambu-Goto modes irrespective of the background conformal frame. Given that the energy reflection/transmission properties of the dual interface are tunable, it would be of great importance to study how (36) is realized for the dual interface both in case of the static holographic junction and the holographic junction corresponding to a nontrivial solution of the nonlinear Nambu-Goto equation when the full nonlinear gravitational perturbations are included.

Conclusions. In this Letter, we have demonstrated that the gravitational two-way junction including its stringy degrees of freedom can be decoded in terms of quantum maps at the level of linearized gravitational perturbations. In the absence of stringy vibrations, these maps reduce to frequency independent universal maps [20,21,36,37] that are realized by a defect operator, while the stringy modes correspond to generalized maps involving compositions with relative conformal transformations. We have particularly improved the methodology of [33] to show that the transmission and reflection on both sides of the interface occur independently, and the full quantum maps are independent of the conformal frame as can be anticipated from the literature [20]. Our results imply generalization of the conformal interfaces studied in [20,21,36,37] realizing tunable energy transmitters. We have argued that the generalized quantum maps are relevant even for the static holographic interface when the full nonlinear gravitational perturbations are included.

It would be also of profound interest to understand how explicit reconstruction of subregions of the gravitational junction from interface CFT subregions can work by generalizing the reformulation of bulk reconstruction in holography in terms of recovery maps of quantum error correcting codes [4,7].

It would also be interesting to investigate the quantum null energy condition [38–40] in holographic interfaces dual to gravitational junctions. Understanding how the link between quantum energy conditions and quantum thermodynamic bounds demonstrated in [41–44] (implying generalized Clausius inequalities for instance) are realized in interfaces is fundamental for actualization of applicable quantum devices, tunable transmitters for instance, via interfaces in many-body systems.

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