Holographic de Sitter Geometry from Entanglement in Conformal Field Theory

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Explicit Examples of States.— Local propagation of $\delta S$ in auxiliary de Sitter geometry holds for any states which are small perturbations of the vacuum state everywhere in space. In particular, we require that for the excitations under consideration, the first law applies not just for a particular sphere but for all spheres (and their complements) on a given time slice. Below we discuss two examples of such states.

Let us first consider a pure state in a $d$-dimensional CFT on a plane, which is created by an infinitesimal insertion of the energy density operator $T_{tt}$ at time $t_0 + i \tau$ and position $x_0$

$$|\phi\rangle = (1 + \epsilon T_{tt})|0\rangle,$$

see also Ref. [19]. The parameter $\epsilon$ is taken to be small in the sense of $\epsilon/\tau^d \ll 1$. The evolution in imaginary time $\tau$ is included to regulate potential UV divergences and ensures that the state is a small perturbation of the vacuum.

The energy density of the state $|S_1\rangle$ is determined by the two-point function of the stress tensor $S_{\mu\nu}$.

$$\langle \phi|T_{tt}(t,x)|\phi\rangle = \epsilon C_T \left[ \frac{1}{(\Delta x^2 - (\Delta t + i \tau)^2)^d} \right.\times \left. \left( \frac{(\Delta x^2 + (\Delta t + i \tau)^2)^2}{(\Delta x^2 - (\Delta t + i \tau)^2)^2} - \frac{1}{d} \right) + c.c. \right] + \mathcal{O}(\epsilon^2),$$

where $C_T$ is the central charge and we have defined $\Delta x^2 = |\vec{x} - \vec{x}_0|^2$ and $\Delta t^2 = |t - t_0|^2$. An explicit illustration of the energy density for this state and corresponding dS propagation of the perturbation in the EE is shown in Fig. S1. Note that the energy density $|S_2\rangle$ is a spherical shell expanding out from $(t_0, \vec{x}_0)$ at the speed of light. As expected from our general argument, the energy density profile $|S_2\rangle$ obeys the constraints (15) and, hence, the holographic propagation respects the antipodal symmetry on the auxiliary dS background.

Our second example is the following mixed state

$$\rho = |0\rangle\langle 0| + \eta |E\rangle\langle E|,$$

where $|E\rangle$ is an energy eigenstate (with constant energy density), and $\eta$ is a small parameter. In this case, we assume that the constant time slice has topology $S^{d-1}$ with radius $r$. Let us now look at $(d-2)$-dimensional spherical entangling surfaces surrounding a cap of the $S^{d-1}$ specified by the angle $\theta_0$. The first law reads now

$$\delta S = 2\pi \int_0^{\theta_0} \frac{\sin^{d-2}\theta}{r^{d-1}\Omega_{d-1}} \times \Omega_{d-2} \cos \theta - \cos \theta_0 \times \frac{\eta E}{\sin \theta_0} \times \frac{1}{r^{d-1}\Omega_{d-1}}\times,$$

where $\Omega_n = 2\pi^{(n+1)/2}/\Gamma(\frac{n+1}{2})$ denotes the volume of a unit $S^n$. The factors in the integrand then correspond to, in order, the volume element of $S^{d-1}$, the boundary-to-bulk propagator for $dS$ in global coordinates, and the (constant) expectation value of the energy density. A special case of this expression appears in Ref. [22], which discusses universal thermal corrections to the vacuum entanglement entropy. There, the energy is given by $E = \frac{\beta}{\Delta}$, where $\Delta$ is the smallest scaling dimension in the spectrum (apart from the identity), and $\eta$ is the product of the degeneracy of the energy eigenstate and the corresponding Boltzmann factor, i.e., $\eta = g e^{-\beta\Delta/r}$. Clearly, this and other mixed states of the form $|S_3\rangle$ violate the first constraint in Eq. (17). Hence, the corresponding $\delta S$ propagates on $dS$ without antipodal symmetry.

Note that $\delta S$ in Eq. (S4) diverges as $\theta_0 \to \pi$, i.e., as the $dS$ propagation reaches the past boundary $T^-$. This divergence is related to a breakdown of the first law and corresponding free propagation in dS space when $\sin \theta_0 \sim \eta E r$ (with $\theta_0 > \pi/2$).

Alternative derivation of wave equation on dS geometry.— Let us now present another perspective on the wave equation (3). The conformal group relevant for a $d$-dimensional CFT is $SO(2,d)$. However, only the subgroup $SO(1,d)$ leaves a constant time slice invariant. Hence the corresponding spherical entangling surfaces are
mapped onto one another under the action of SO(1, d).
Now considering the perturbations $\delta S$ for these ball-
shaped regions, the SO(1, d) generators $K_i$ act as

$$\partial K_i \delta S \left[ \langle T_{tt} \rangle \right] = -\delta S \left[ \langle \partial K_i T_{tt} \rangle \right].$$

(55)

Here, the $\partial K_i$ on the RHS can be viewed as generat-
ing an “active” conformal transformation that changes
the CFT state, while the $\partial K_i$ on the LHS generates a
“passive” transformation that instead changes the spher-
ical entangling surface. Comparing now the “active” and
“passive” action of the quadratic Casimir of SO(1, d),
$\nabla^2 \equiv c_{ij} \partial K_i \partial K_j$, we get

$$\nabla^2 \delta S \left[ \langle T_{tt} \rangle \right] = \delta S \left[ \langle \nabla^2 T_{tt} \rangle \right] = -d \delta S \left[ \langle T_{tt} \rangle \right],$$

(56)

where the second expression above uses the linearity
in $T_{tt}$ of the modular Hamiltonian (8). Further, the
last expression appears after using the fact that the en-
ergy density transforms as a scalar of weight $d$ with
respect to the SO(1, d) subgroup. Now, a particu-
lar spherical entangling surface is left invariant by the
stabilizer group SO(1, d − 1). Hence, on the LHS of
Eq. (56), the nontrivial action of $\nabla^2$ is on the coset space
SO(1, d)/SO(1, d − 1). The latter coset is precisely the
anticipated $d$-dimensional dS geometry, and $\nabla^2$ becomes
the d’Alembertian on this space. Hence this group theo-
retic approach produces precisely the Klein-Gordon equa-
tion (3) on the auxiliary dS space. Note that this analysis
implicitly normalizes the dS radius $L$ to unity.

Finally, let us mention that the group theoretic argu-
ment above can be also generalized to the higher-spin
case and, as expected, yields the mass given by Eq. (17).


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