

This is the accepted manuscript made available via CHORUS. The article has been published as:

# Diagnosing Chaos Using Four-Point Functions in Two-Dimensional Conformal Field Theory

Daniel A. Roberts and Douglas Stanford

Phys. Rev. Lett. **115**, 131603 — Published 22 September 2015

DOI: [10.1103/PhysRevLett.115.131603](https://doi.org/10.1103/PhysRevLett.115.131603)

# Diagnosing chaos using four-point functions in two-dimensional conformal field theory

Daniel A. Roberts\*

*Center for Theoretical Physics and Department of Physics,  
Massachusetts Institute of Technology, Cambridge, MA, USA*

Douglas Stanford†

*School of Natural Sciences, Institute for Advanced Study, Princeton, NJ, USA*

We study chaotic dynamics in two-dimensional conformal field theory through out-of-time-order thermal correlators of the form  $\langle W(t) V W(t) V \rangle$ . We reproduce holographic calculations similar to those of [1], by studying the large  $c$  Virasoro identity conformal block. The contribution of this block to the above correlation function begins to decrease exponentially after a delay of  $\sim t_* - \frac{\beta}{2\pi} \log \beta^2 E_w E_v$ , where  $t_*$  is the fast scrambling time  $\frac{\beta}{2\pi} \log c$ , and  $E_w, E_v$  are the energy scales of the  $W, V$  operators.

---

\* drob@mit.edu

† stanford@ias.edu

## INTRODUCTION

In classical mechanics, the butterfly effect is a vivid diagnostic of chaos: small perturbations grow rapidly to affect the entire system. In quantum mechanics, studies of chaos have often focused on less direct measures, such as the statistics of level spacings in the energy spectrum [2, 3], or the local properties of energy eigenvectors [4–6]. In this paper, following [1, 7–9] we will focus on a notion of quantum chaos that is closely related to the classical butterfly effect.

Consider a pair of rather general Hermitian operators  $V$  and  $W$  in a quantum mechanical system. If the system is strongly chaotic, we expect perturbations by  $V$  to affect later measurements of  $W$ , for almost any choice of operators  $V$  and  $W$  [10]. In other words, we expect the commutator  $[V, W(t)]$  to become large. A useful diagnostic is the expectation value of the square of the commutator [11]:

$$-\langle [V, W(t)]^2 \rangle_\beta = \langle VW(t)W(t)V \rangle_\beta + \langle W(t)VVW(t) \rangle_\beta \quad (1)$$

$$- \langle VW(t)VW(t) \rangle_\beta - \langle W(t)VW(t)V \rangle_\beta. \quad (2)$$

Here  $\langle \cdot \rangle_\beta$  indicates the thermal expectation value at inverse temperature  $\beta$ . Let us assume that the operators commute at  $t = 0$ . For small times, the terms on the first line cancel the terms on the second line. As we move to large time, the terms on the first line will each approach the order-one value  $\langle WW \rangle_\beta \langle VV \rangle_\beta$ . This can be understood by viewing the  $VWWV$  ordering as an expectation value of  $WW$  in a state given by acting with  $V$  on the thermal state. If the energy injected by  $V$  is small, the state will relax and the expectation value will approach the thermal value,  $\langle WW \rangle_\beta$ , multiplied by the norm of the state,  $\langle VV \rangle_\beta$ .

By contrast, in a suitably chaotic system, the correlation functions on the second line will become small for large  $t$ . From (1), it is clear that this will imply a large commutator, and hence a quantum butterfly effect. We believe that this happens for practically any [12] choice of  $W$  and  $V$  and that, in fact, this behavior is a basic diagnostic of quantum chaos.

This behavior has been confirmed for theories holographically dual to Einstein gravity in recent work [1, 7–9, 13]. There, the thermal state is represented by a black hole, and the  $V, W$  operators create quanta that collide near the horizon. The key effect leading to a large commutator is the exponential blueshift relative to later slices as the  $W$  perturbation falls into the black hole.

The purpose of this letter is to reproduce part of that analysis without using holography directly. We will work in a 2d conformal field theory (CFT) where thermal expectation values are related to vacuum expectation values by a conformal transformation. We will also restrict attention to a particular contribution to the four point function given by the large  $c$  Virasoro identity conformal block. This resums the terms corresponding to factorization on powers and derivatives of the stress tensor. The close relationship between 2+1 gravity and the identity Virasoro block has been demonstrated recently in [14–16]; our work should be understood as an application of the techniques in these papers to the problem studied in [1, 7–9, 13]. We will find exact agreement between the Virasoro block calculation and the corresponding holographic calculation in pure 3d gravity. In particular, we will find that the contribution of the identity block to  $\langle VW(t)VW(t) \rangle_\beta$  begins to decay exponentially around the fast scrambling time [17, 18]  $t_* = \frac{\beta}{2\pi} \log c$ .

A key point that will emerge in our analysis is the following. Each of the Lorentzian correlators on the right hand side of (1) can be obtained by analytic continuation of the same Euclidean four point function. The sharp difference in behavior between the first line and the second line arises because the continuation defines a multivalued function, with different orderings corresponding to different sheets. To see the butterfly effect (via the decay of the terms on the second line), one has to move off the principal sheet.

It is important to emphasize that although the Virasoro identity block does reproduce the gravitational calculation of the four point function, it is not the full answer. Indeed, a single Virasoro primary with spin greater than two can easily dominate the contribution from the identity in a certain range of times  $t$ . In holographic terms, this sensitivity to the spectrum is related to the fact that the high energy collision depends on stringy corrections. Just as the gravitational calculation is a useful model for a more accurate

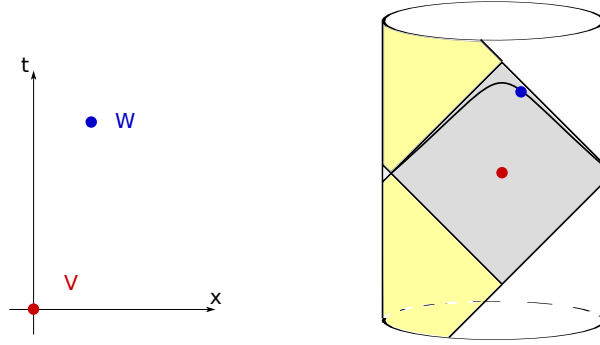


FIG. 1. **Left:** the spacetime arrangement of the  $W$  and  $V$  operators. **Right:** their locations after the conformal mapping, viewed in the Rindler patch on the boundary of  $\text{AdS}_3$  (grey) covered by  $x, t$ . The union of the grey and yellow regions are the Poincare patch covered by  $z, \bar{z}$ .

string-corrected analysis [19], the Virasoro identity block provides a model for a more complete CFT calculation.

Although it is not our purpose to prove that the four-point-function diagnostic described above agrees with other definitions of quantum chaos, we will provide a sanity check by evaluating the above four-point functions in the two-dimensional Ising model. For this system, we will see that certain out-of-time-order four-point functions do not tend to zero for large  $t$ .

## CFT CALCULATIONS

### Conventions and review

In this paper, we will study thermal four-point correlation functions of  $W$  and  $V$  of the form in (1) and (2). Eventually, these operators will be arranged in the timelike configuration shown in Fig. 1, where  $V$  is at the origin, and  $W$  is at position  $t > x > 0$ . However, we will obtain these correlation functions by starting with the Euclidean correlator and analytically continuing. In 2d CFT, we can map thermal expectation values to vacuum expectation values through the conformal transformation

$$z(x, t) = e^{\frac{2\pi}{\beta}(x+t)}, \quad \bar{z}(x, t) = e^{\frac{2\pi}{\beta}(x-t)}. \quad (3)$$

Here,  $x, t$  are the original coordinates on the spatially infinite thermal system and  $z, \bar{z}$  are coordinates on the vacuum system. Explicitly,

$$\langle \mathcal{O}(x, t) \dots \rangle_{\beta} = \left( \frac{2\pi z}{\beta} \right)^h \left( \frac{2\pi \bar{z}}{\beta} \right)^{\bar{h}} \langle \mathcal{O}(z, \bar{z}) \dots \rangle, \quad (4)$$

where  $h, \bar{h}$  are the conformal weights of the  $\mathcal{O}$  operator, related to the dimension and spin by  $\Delta = h + \bar{h}$  and  $J = h - \bar{h}$ . On the left hand side, we have a thermal expectation value, at inverse temperature  $\beta$ , and on the right hand side we have a vacuum expectation value on the  $z, \bar{z}$  space. It is common to work with units in which  $\beta = 2\pi$ , but we prefer to keep the  $\beta$  dependence explicit.

It will be essential in this paper to study correlation functions with operators at complexified times  $t_i$ . In our convention, real  $t$  corresponds to Minkowski time, and imaginary  $t$  corresponds to Euclidean time. Notice from (3) that  $\bar{z}$  is the complex conjugate of  $z$  only if the time  $t$  is purely Euclidean. In order to make contact with standard CFT formulas for the four point function, we will begin with a purely Euclidean arrangement of the operators. This means a choice of  $z_1, \bar{z}_1, \dots, z_4, \bar{z}_4$  with  $\bar{z}_i = z_i^*$ . With such

a configuration, the ordering of the operators is unimportant, and global conformal invariance on the  $z, \bar{z}$  plane implies that the four point function can be written

$$\langle W(z_1, \bar{z}_1) W(z_2, \bar{z}_2) V(z_3, \bar{z}_3) V(z_4, \bar{z}_4) \rangle = \frac{1}{z_{12}^{2h_w} z_{34}^{2h_v}} \frac{1}{\bar{z}_{12}^{2\bar{h}_w} \bar{z}_{34}^{2\bar{h}_v}} f(z, \bar{z}), \quad (5)$$

in terms of a function  $f$  of the conformally invariant cross ratios

$$z = \frac{z_{12} z_{34}}{z_{13} z_{24}}, \quad \bar{z} = \frac{\bar{z}_{12} \bar{z}_{34}}{\bar{z}_{13} \bar{z}_{24}}, \quad z_{ij} \equiv z_i - z_j. \quad (6)$$

According to the general principles of CFT, we can expand  $f$  as a sum of global conformal blocks, explicitly [20, 21]

$$f(z, \bar{z}) = \sum_{h, \bar{h}} p(h, \bar{h}) z^h \bar{z}^{\bar{h}} F(h, h, 2h, z) F(\bar{h}, \bar{h}, 2\bar{h}, \bar{z}) \quad (7)$$

where  $F$  is the Gauss hypergeometric function, the sum is over the dimensions of global  $SL(2)$  primary operators, and the constants  $p$  are related to OPE coefficients  $p(h, \bar{h}) = \lambda_{WW\mathcal{O}_{h, \bar{h}}} \lambda_{VV\mathcal{O}_{h, \bar{h}}}$ .

### Continuation to the second sheet

In order to apply the above formulas to the correlators (1) and (2), we need to understand how to obtain them as analytic continuations of the Euclidean four point function. That this is possible follows from the fact that all Wightman functions are analytic continuations of each other [22]. The procedure involves three steps: first, one starts with the Euclidean function, assigning small and different imaginary times  $t_j = i\epsilon_j$  to each of the operators. Second, with the imaginary times held fixed, one increases the real times of the operators to the desired Lorentzian values. Finally, one smears the operators in real time and then takes the imaginary times  $\{\epsilon_i\}$  to zero [23]. The result will be a Lorentzian correlator ordered such that the leftmost operator corresponds to the smallest value of  $\epsilon$ , the second operator corresponds to the second smallest, and so on.

This elaborate procedure is necessary because Eq. (5) is a multivalued function of the independent complex variables  $\{z_i, \bar{z}_i\}$ . The interesting multivaluedness comes from  $f(z, \bar{z})$ . By crossing symmetry, this function is single valued on the Euclidean section  $\bar{z} = z^*$ , but it is multivalued as a function of independent  $z$  and  $\bar{z}$ , with branch cuts extending from one to infinity. Different orderings of the  $W, V$  operators correspond to different sheets of this function. To determine the correct sheet, we must assign  $i\epsilon$ 's as above, and follow the path of the cross ratios, watching to see if they pass around the branch loci at  $z = 1$  and  $\bar{z} = 1$ .

To carry this out directly, we write

$$z_1 = e^{\frac{2\pi}{\beta}(t' + i\epsilon_1)} \quad \bar{z}_1 = e^{-\frac{2\pi}{\beta}(t' + i\epsilon_1)} \quad (8)$$

$$z_2 = e^{\frac{2\pi}{\beta}(t' + i\epsilon_2)} \quad \bar{z}_2 = e^{-\frac{2\pi}{\beta}(t' + i\epsilon_2)} \quad (9)$$

$$z_3 = e^{\frac{2\pi}{\beta}(x + i\epsilon_3)} \quad \bar{z}_3 = e^{\frac{2\pi}{\beta}(x - i\epsilon_3)} \quad (10)$$

$$z_4 = e^{\frac{2\pi}{\beta}(x + i\epsilon_4)} \quad \bar{z}_4 = e^{\frac{2\pi}{\beta}(x - i\epsilon_4)} \quad (11)$$

as a function of the continuation parameter  $t'$ . When  $t' = 0$ , we have a purely Euclidean correlator, on the principal sheet of the function  $f(z, \bar{z})$ . When  $t' = t > x$ , we have an arrangement of operators as shown in Fig. 1.

The cross ratios  $z, \bar{z}$  are determined by these coordinates as in Eq. (6). Their paths, as a function of  $t'$ , depend on the ordering of operators through the associated  $i\epsilon$  prescription. Representative paths for the

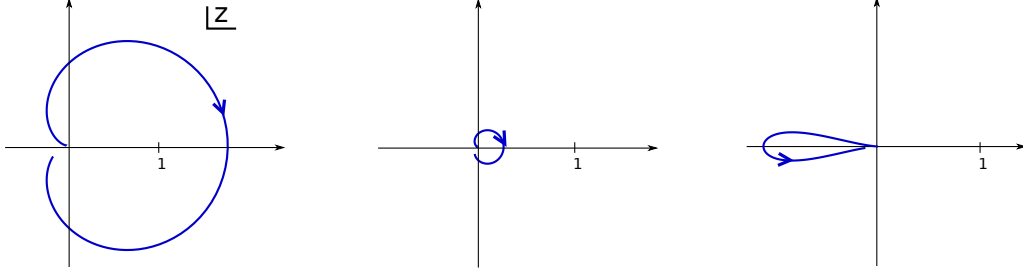


FIG. 2. The paths taken by the cross ratio  $z$  during the continuations corresponding to (from left to right)  $\langle WVWV \rangle$ ,  $\langle WWVV \rangle$ , and  $\langle WVWV \rangle$ . Only in the first case does the path pass around the branch point at  $z = 1$ .

three cases of interest are shown in Fig. 2. The variable  $\bar{z}$  never passes around the branch point at one, and the  $z$  variable does so only in the case corresponding to  $WVWV$  [24].

In the final configuration with  $t' = t$ , the cross ratios are small. For  $t \gg x$ , we have

$$z \approx -e^{\frac{2\pi}{\beta}(x-t)} \epsilon_{12}^* \epsilon_{34} \quad \bar{z} \approx -e^{-\frac{2\pi}{\beta}(x+t)} \epsilon_{12}^* \epsilon_{34}, \quad (12)$$

where we introduced the abbreviation

$$\epsilon_{ij} = i(e^{\frac{2\pi}{\beta}i\epsilon_i} - e^{\frac{2\pi}{\beta}i\epsilon_j}). \quad (13)$$

For the orderings  $WWVV$  and  $WVWV$ , no branch cuts are crossed, so the limit of small cross ratios can be taken on the principal sheet of Eq. (7). The contribution from the identity operator dominates, verifying our statement in the Introduction that both  $\langle W(t)VWV(t) \rangle_\beta$  and  $\langle W(t)W(t)VWV(t) \rangle_\beta$  approach  $\langle WW \rangle \langle VV \rangle_\beta$  for large  $t$ .

For  $WVWV$ ,  $z$  passes around the branch point at one. The hypergeometric function  $F(a, b, c, z)$  has known monodromy around  $z = 1$ , returning to a multiple of itself, plus a multiple of the other linearly independent solution to the hypergeometric equation,  $z^{1-c}F(1+a-c, 1+b-c, 2-c, z)$ . For small  $z, \bar{z}$  we then have

$$f(z, \bar{z}) \approx \sum_{h, \bar{h}} \tilde{p}(h, \bar{h}) z^{1-h} \bar{z}^{\bar{h}}, \quad (14)$$

where  $\tilde{p}$  has been defined to absorb the transformation coefficient. On this sheet, as  $z, \bar{z}$  become small, global primaries with large spin become important. As a function of  $x, t$ , individual terms in this sum grow like  $e^{(h-\bar{h}-1)t} e^{-(h+\bar{h}-1)x}$ . For sufficiently large  $t$ , this sum diverges, and it must be defined by analytic continuation. In other words, we must do the sum over  $h, \bar{h}$  before we continue the cross ratios. In a CFT dual to string theory in  $\text{AdS}_3$ , we expect this divergence even at a fixed order in the large  $c$  expansion, because of the sum over higher spin bulk exchanges [19].

### The Virasoro identity block

The primary focus of this paper is reproducing the Einstein gravity calculation of the correlation function. This calculation was done by studying free propagation on a shock wave background, which implicitly sums an infinite tower of ladder exchange diagrams. In the CFT, these diagrams are related to terms involving powers and derivatives of the stress tensor in the OPE representation of the four point function. In a two-dimensional CFT, all such terms can be treated simultaneously using the Virasoro conformal block of the identity operator, itself an infinite sum of  $SL(2)$  conformal blocks. Including only these terms in the OPE amounts to replacing

$$f(z, \bar{z}) \rightarrow \mathcal{F}(z) \bar{\mathcal{F}}(\bar{z}). \quad (15)$$

where  $\mathcal{F}$  is the Virasoro conformal block with dimension zero in the intermediate channel. This substitution is appropriate for a large  $N$  CFT with a sparse spectrum of single trace higher spin operators [25].

The function  $\mathcal{F}$  is not known explicitly, but there are several methods for approximating it [14–16]. We will use a formula from [15], which is valid at large  $c$ , with  $h_w/c$  fixed and small, and  $h_v$  fixed and large. Here, the formula reads

$$\mathcal{F}(z) \approx \left( \frac{z(1-z)^{-6h_w/c}}{1-(1-z)^{1-12h_w/c}} \right)^{2h_v}. \quad (16)$$

This function has a branch point at  $z = 1$ , as expected. Following the contour around  $z = 1$  and taking  $z$  small, we find

$$\mathcal{F}(z) \approx \left( \frac{1}{1 - \frac{24\pi i h_w}{cz}} \right)^{2h_v}. \quad (17)$$

The trajectory of  $\bar{z}$  does not circle the branch point at  $\bar{z} = 1$ , so for small  $\bar{z}$ , we simply have  $\bar{\mathcal{F}}(\bar{z}) \approx 1$ , the contribution of the identity operator itself. Substituting (17) in (15) and then in (5), we find

$$\frac{\langle W(t + i\epsilon_1)V(i\epsilon_3)W(t + i\epsilon_2)V(i\epsilon_4) \rangle_\beta}{\langle W(i\epsilon_1)W(i\epsilon_2) \rangle_\beta \langle V(i\epsilon_3)V(i\epsilon_4) \rangle_\beta} \approx \left( \frac{1}{1 + \frac{24\pi i h_w}{\epsilon_{12}^* \epsilon_{34}} e^{\frac{2\pi}{\beta}(t-t_*-x)}} \right)^{2h_v} \quad (18)$$

where we define the fast scrambling time  $t_*$  [17, 18] with the convention

$$t_* = \frac{\beta}{2\pi} \log c. \quad (19)$$

Eq. (18) is the main result of our paper. The correlation function (18) begins to decrease at a time  $t \sim \frac{\beta}{2\pi} \log c/h_w$ .

This formula agrees precisely with the bulk analysis (reviewed in the Appendix) in the above scaling. It is also interesting to consider the scaling  $h_v, h_w$  fixed,  $h_w \gg h_v \gg 1$  with  $c \rightarrow \infty$ . The bulk analysis suggests that Eq. (18) is also correct in this scaling, but even without the Virasoro block analysis, the  $SL(2)$  block of the stress tensor (which gives a contribution  $\propto h_w h_v/(cz)$  on the second sheet) is enough to show that for general dimensions, the time until the identity Virasoro block is affected is of order  $t_* - \frac{\beta}{2\pi} \log h_w h_v$ , where the second term is order one in this scaling. This gives a field-theoretic explanation for the origin of the fast scrambling time.

We are grateful to Ethan Dyer, Tom Hartman, Alexei Kitaev, and Steve Shenker for helpful discussions. D.R. is supported by the Fannie and John Hertz Foundation and is very thankful for the hospitality of the Stanford Institute for Theoretical Physics during the completion of this work. D.R. also acknowledges the U.S. Department of Energy under cooperative research agreement Contract Number DE-SC00012567. D.S. is supported by NSF grant PHY\_1314311/Dirac.

### Bulk calculations

In this section, we will compute the correlation function (18) using the gravitational shock wave methods of [1, 8]. If  $h_w \gg h_v \gg 1$ , we can calculate the correlation function by treating the  $W$  operator as creating a shock wave, and calculating the two point function of the  $V$  operator on that background. The analysis breaks into two parts: (i) finding the geometry of the shock sourced by  $W$  and (ii) computing the correlation function of the  $V$  operators in that background.

The metric of a localized shock wave [26] in 2+1 dimensional AdS-Rindler space, is [1, 8, 27–29]

$$ds^2 = -\frac{4}{(1+uv)^2} du dv + \frac{(1-uv)^2}{(1+uv)^2} dx^2 + 4\delta(u)h(x)du^2, \quad (20)$$

where  $h$  will be defined below. We will consider a shock sourced by a stress tensor

$$T_{uu}(u', v', x') = P \delta(u') \delta(x' - x), \quad (21)$$

appropriate for a particle sourced by the  $W$  operator, traveling along the  $u = 0$  horizon at transverse position  $x$ . The metric (20) can be understood as two halves of AdS-Rindler, glued together at  $u = 0$  with a shift

$$\delta v(x) = h(x) \quad (22)$$

in the  $v$  direction. Plugging in to Einstein's equations, we determine  $h$  as

$$h(x') = 2\pi G_N P e^{-|x' - x|}. \quad (23)$$

To relate this geometry to the state  $|\psi\rangle$ , we need to fix  $P$ . In other words, we need to evaluate

$$\frac{\langle \psi | \int dx' du' T_{uu} | \psi \rangle}{\langle \psi | \psi \rangle}, \quad (24)$$

where we take the integral to run over the slice  $v = 0$ . We will assume that  $W$  is dual to a single-particle operator in the bulk, so that the state  $|\psi\rangle$  can be described by a Klein-Gordon wave function  $K$ . This wave function is a bulk-to boundary propagator from the location  $(x, t)$  of the  $W$  operator. It is given in terms of the regularized geodesic distance  $d$  from the boundary point, as  $K \propto (\cosh d)^{-2h_w}$ . At  $v = 0$ , we find

$$K(t, x; u', x') = \frac{\mathcal{N}}{[e^t u' + \cosh(x - x')]^{2h_w}}. \quad (25)$$

The norm  $\langle \psi | \psi \rangle$  is a Klein-Gordon inner product

$$\langle \psi | \psi \rangle = 2i \int dx' du' K(t + i\tau, x; u', x')^* \partial_{u'} K(t + i\tau, x; u', x'). \quad (26)$$

The  $u'$  integral can be done using contour integration, and the  $x'$  integral can be done in terms of  $\Gamma$  functions:

$$\langle \psi | \psi \rangle = \mathcal{N}^2 \frac{4\pi^{3/2}}{(2 \sin \tau)^{4h_w}} \frac{\Gamma[4h_w]}{\Gamma[2h_w] \Gamma[2h_w + \frac{1}{2}]}. \quad (27)$$

For the numerator, the stress tensor for the Klein-Gordon field is given by the expression  $T_{uu} = \partial_u \varphi \partial_u \varphi$ . Contracting bulk operators with boundary operators using  $K$ , we have

$$\langle \psi | \int dx' du' T_{uu} | \psi \rangle = 2 \int dx' du' \partial_{u'} K(t_w + i\tau, x; u', x')^* \partial_{u'} K(t_w + i\tau, x; u', x'), \quad (28)$$

where the factor of 2 comes from the two different ways of doing the contractions. The integrals can be done the same way as before:

$$\langle \psi | \int dx' du' T_{uu} | \psi \rangle = \mathcal{N}^2 \frac{8\pi^{3/2} e^{t_w}}{(2 \sin \tau)^{4h_w+1}} \frac{\Gamma[4h_w] \Gamma[2h_w + \frac{1}{2}]}{\Gamma[2h_w]^3}. \quad (29)$$

Taking the ratio at large  $h_w$ , we find

$$P = \frac{2h_w e^{t_w}}{\sin \tau}. \quad (30)$$



The second step, following [1], is to compute the two-sided correlation function of the  $V$  operators in this shock background. We will do this using the geodesic approximation

$$\langle \psi | V_L V_R | \psi \rangle \propto e^{-md}, \quad (31)$$

where  $d$  is the regularized geodesic distance, and the mass  $m$  is approximately  $2h_v$ , the conformal weight of  $V$ . Following [1], we find

$$d = 2 \log 2r_\infty + \log [1 + h(0)]. \quad (32)$$

After subtracting the divergent distance in the unperturbed thermofield double state,  $d_{TFD} = 2 \log 2r_\infty$ , we plug the distance into (31), finding

$$\frac{\langle \psi | V_L V_R | \psi \rangle}{\langle \psi | \psi \rangle \langle V_L V_R \rangle} = \left( \frac{1}{1 + \frac{4\pi G_N h_w}{\sin \tau} e^{t-x}} \right)^{2h_v}. \quad (33)$$

This agrees with (18) after (i) using  $G_N = 3/2c$  to express the gravitational constant in terms of the central charge, (ii) plugging in  $\epsilon_1 = -\tau$ ,  $\epsilon_2 = \tau$ ,  $\epsilon_3 = 0$ ,  $\epsilon_4 = \beta/2$  and (iii) using  $\beta = 2\pi$ .

### Ising model

It is interesting to contrast the behavior of out-of-time-order correlators in a chaotic theory to those in an integrable theory. As an example, we consider the two-dimensional Ising CFT. This theory has  $c = 1/2$  and three Virasoro primary operators:  $I$ ,  $\sigma$ , and  $\epsilon$ , corresponding to the identity, ‘spin,’ and ‘energy’ operators. The different combinations of four-point correlators of these primaries are well known [30–32]. We will present these by giving the functions  $f(z, \bar{z})$  in Eq.(5):

$$f_{\sigma\sigma}(z, \bar{z}) = \frac{1}{2} \left( |1 + \sqrt{1-z}| + |1 - \sqrt{1-z}| \right), \quad (34)$$

$$f_{\sigma\epsilon}(z, \bar{z}) = \left| \frac{2-z}{2\sqrt{1-z}} \right|^2, \quad (35)$$

$$f_{\epsilon\epsilon}(z, \bar{z}) = \left| \frac{1-z+z^2}{1-z} \right|^2, \quad (36)$$

where the operators are ordered  $WVWV$  with the configuration specified by (8)–(11). Following the contour across the branch cut for the two correlators that do have a second sheet and taking  $z$  small, we find

$$\frac{\langle \sigma\sigma\sigma\sigma \rangle_\beta}{\langle \sigma\sigma \rangle_\beta^2} = 0, \quad \frac{\langle \sigma\epsilon\sigma\epsilon \rangle_\beta}{\langle \sigma\sigma \rangle_\beta \langle \epsilon\epsilon \rangle_\beta} = -1, \quad \frac{\langle \epsilon\epsilon\epsilon\epsilon \rangle_\beta}{\langle \epsilon\epsilon \rangle_\beta^2} = 1. \quad (37)$$

Only  $\langle \sigma\sigma\sigma\sigma \rangle$  vanishes at large  $t$ .

- 
- [1] S. H. Shenker and D. Stanford, “Black holes and the butterfly effect,” [arXiv:1306.0622 \[hep-th\]](#).
  - [2] M. V. Berry and M. Tabor, “Level clustering in the regular spectrum,” *Proceedings of the Royal Society of London. A. Mathematical and Physical Sciences* **356** no. 1686, (1977) 375–394.
  - [3] O. Bohigas, M. J. Giannoni, and C. Schmit, “Characterization of chaotic quantum spectra and universality of level fluctuation laws,” *Phys. Rev. Lett.* **52** (Jan, 1984) 1–4. <http://link.aps.org/doi/10.1103/PhysRevLett.52.1>.

- [4] M. Berry, “Regular and irregular semiclassical wavefunctions,” *Journal of Physics A: Mathematical and General* **10** no. 12, (1977) 2083.
- [5] J. M. Deutsch, “Quantum statistical mechanics in a closed system,” *Physical Review A* **43** no. 4, (1991) 2046.
- [6] M. Srednicki, “Chaos and quantum thermalization,” *Physical Review E* **50** no. 2, (1994) 888.
- [7] S. H. Shenker and D. Stanford, “Multiple Shocks,” [arXiv:1312.3296 \[hep-th\]](#).
- [8] D. A. Roberts, D. Stanford, and L. Susskind, “Localized shocks,” [arXiv:1409.8180 \[hep-th\]](#).
- [9] A. Kitaev, “Hidden correlations in the hawking radiation and thermal noise.”. Talk given at the Fundamental Physics Prize Symposium, Nov. 10, 2014.
- [10] We take  $V, W$  to be approximately local operators, smeared over a thermal scale, and with one point functions subtracted.
- [11] A. Larkin and Y. Ovchinnikov, “Quasiclassical method in the theory of superconductivity,” *JETP* **28,6** (1969) 1200–1205.
- [12] Here, “practically any” should include all local operators with both  $h, \bar{h}$  nonzero.
- [13] S. Leichenauer, “Disrupting Entanglement of Black Holes,” *Phys.Rev.* **D90** (2014) 046009, [arXiv:1405.7365 \[hep-th\]](#).
- [14] T. Hartman, “Entanglement Entropy at Large Central Charge,” [arXiv:1303.6955 \[hep-th\]](#).
- [15] A. L. Fitzpatrick, J. Kaplan, and M. T. Walters, “Universality of Long-Distance AdS Physics from the CFT Bootstrap,” *JHEP* **1408** (2014) 145, [arXiv:1403.6829 \[hep-th\]](#).
- [16] C. T. Asplund, A. Bernamonti, F. Galli, and T. Hartman, “Holographic Entanglement Entropy from 2d CFT: Heavy States and Local Quenches,” [arXiv:1410.1392 \[hep-th\]](#).
- [17] P. Hayden and J. Preskill, “Black holes as mirrors: Quantum information in random subsystems,” *JHEP* **0709** (2007) 120, [arXiv:0708.4025 \[hep-th\]](#).
- [18] Y. Sekino and L. Susskind, “Fast Scramblers,” *JHEP* **0810** (2008) 065, [arXiv:0808.2096 \[hep-th\]](#).
- [19] S. H. Shenker and D. Stanford, “Stringy effects in scrambling,” [arXiv:1412.6087 \[hep-th\]](#).
- [20] A. Zamolodchikov, “Conformal symmetry in two-dimensions: An explicit recurrence formula for the conformal partial wave amplitude,” *Commun.Math.Phys.* **96** (1984) 419–422.
- [21] F. Dolan and H. Osborn, “Conformal four point functions and the operator product expansion,” *Nucl.Phys.* **B599** (2001) 459–496, [arXiv:hep-th/0011040 \[hep-th\]](#).
- [22] Theorem 3.6 of [33].
- [23] In fact, we will omit this final step in this paper. However, we will omit it consistently on both sides of the bulk and boundary calculations that we are comparing.
- [24] A very similar continuation was discussed for high-energy scattering kinematics in [29, 34, 35].
- [25] It is interesting to note that, because of the rapid  $z^{1-h}\bar{z}^{\bar{h}}$  behavior on the second sheet, even one single-trace operator with spin greater than two and dimension parametrically independent of  $c$  will dominate over the Virasoro identity block near time  $t_*$ . This means that the universal contribution to the correlator identified in this paper applies only to CFTs with a very sparse spectrum of single trace operators. This matches the sensitivity of high energy scattering in gravity to massive higher spin fields, which grow more rapidly with energy. In a consistent bulk theory of higher spin fields, such as string theory, a Regge-type resummation of an infinite number of stringy operators is necessary. For further discussion of this point in the present context, see [19].
- [26] Although we refer to these geometries as shock waves, the terminology is somewhat misleading in 2+1 dimensions, since the geometry is locally pure  $\text{AdS}_3$  away from the source.
- [27] T. Dray and G. ’t Hooft, “The Gravitational Shock Wave of a Massless Particle,” *Nucl.Phys.* **B253** (1985) 173–188.
- [28] G. T. Horowitz and N. Itzhaki, “Black holes, shock waves, and causality in the AdS / CFT correspondence,” *JHEP* **9902** (1999) 010, [arXiv:hep-th/9901012 \[hep-th\]](#).
- [29] L. Cornalba, M. S. Costa, J. Penedones, and R. Schiappa, “Eikonal Approximation in AdS/CFT: From Shock Waves to Four-Point Functions,” *JHEP* **0708** (2007) 019, [arXiv:hep-th/0611122 \[hep-th\]](#).
- [30] A. Belavin, A. M. Polyakov, and A. Zamolodchikov, “Infinite Conformal Symmetry in Two-Dimensional Quantum Field Theory,” *Nucl.Phys.* **B241** (1984) 333–380.
- [31] M. P. Mattis, “Correlations in Two-dimensional Critical Theories,” *Nucl.Phys.* **B285** (1987) 671.
- [32] P. H. Ginsparg, “Applied Conformal Field Theory,” [arXiv:hep-th/9108028 \[hep-th\]](#).
- [33] R. F. Streater and A. S. Wightman, *PCT, spin and statistics, and all that*. W.A. Benjamin, Inc., 1964.
- [34] L. Cornalba, M. S. Costa, J. Penedones, and R. Schiappa, “Eikonal Approximation in AdS/CFT: Conformal Partial Waves and Finite N Four-Point Functions,” *Nucl.Phys.* **B767** (2007) 327–351, [arXiv:hep-th/0611123 \[hep-th\]](#).
- [35] L. Cornalba, M. S. Costa, and J. Penedones, “Eikonal approximation in AdS/CFT: Resumming the gravitational loop expansion,” *JHEP* **0709** (2007) 037, [arXiv:0707.0120 \[hep-th\]](#).