Universality of Abrupt Holographic Quenches

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We make an analytic investigation of rapid quenches of relevant operators in \(d\)-dimensional holographic CFT's, which admit a dual gravity description. We uncover a universal scaling behaviour in the response of the system, which depends only on the conformal dimension of the quenched operator in the vicinity of the ultraviolet fixed point of the theory. Unless the amplitude of the quench is scaled appropriately, the work done on a system during the quench diverges in the limit of abrupt quenches for operators with dimension \(\frac{d}{2} \leq \Delta < d\).

Quantum quenches have recently become accessible in laboratory experiments \cite{1}, which has initiated much activity by theoretical physicists to understand such systems. Up until now, most analytic work on the topic of relativistic quantum quenches have assumed that the field theory is at weak coupling \cite{2, 3}. The study of quantum quenches at strong coupling is accessible through the gauge/gravity duality \cite{4}. Much related work studying thermalization in the boundary theory was done by studying the gravity dual under the assumption that the non-equilibrium evolution can be approximated by a uniformly evolving spacetime, e.g., \cite{5, 6, 7, 8}. Other approaches study the evolution of a probe on the static spacetime \cite{9}. The approach of numerically evolving the dual gravity theory was initiated in \cite{10}. Further numerical studies of quenches in a variety of holographic systems were presented in \cite{11, 12, 13, 14, 15, 16, 17, 18}.

In \cite{18, 19}, holography was applied to study quenches of the coupling to a relevant scalar operator in the boundary theory. A numerical approach was taken to study the evolution of the dual scalar field in the bulk spacetime. For fast quenches, evidence was found for a universal scaling of the expectation value of the boundary operator. Similar scaling was observed for the change in energy density, pressure and entropy density. However, no analytic understanding of this behaviour was available.

In this Letter, we investigate these holographic quenches analytically, focusing on the work done by the quench. Unlike \cite{18, 19}, in which the coupling was a given analytic function of time, we abruptly (but with some degree of smoothness) switch on this source at \(t = 0\). The coupling is then varied over a finite interval \(\delta t\) and is held constant afterwards. We find that for fast quenches, the essential physics can be extracted by solving the linearized scalar field equation in the asymptotic AdS geometry. Note that our analysis is naturally driven to this regime by the limit \(\delta t \to 0\). In contrast to \cite{18, 19}, we are not a priori limiting our study to a perturbative expansion in the amplitude of the bulk scalar. Our analytic results also cover any spacetime dimension \(d\) for the boundary theory, whereas \cite{18, 19} were limited to \(d = 4\).

Let us describe the quenches in more detail: The coupling in the boundary theory is determined by the leading non-normalizable mode of the bulk scalar \(\phi\). We set this mode to zero before \(t = 0\), vary it in the interval \(0 < t < \delta t\) and hold it fixed afterwards. Because the energy density can only change while the coupling is changing, we are only interested in the response of the scalar field during the timespan \(0 < t < \delta t\). Further, since the response propagates in from the boundary of the spacetime, the field will only be nonzero within the lightcone \(t = \rho\). Hence to determine the work done, we need only solve for the bulk evolution in the triangular region bounded by this lightcone, the surface \(t = \delta t\) and the AdS boundary, as shown in fig. 1. As is also illustrated, as \(\delta t \to 0\), this triangle shrinks to a small region in the asymptotic spacetime. The normalizable component of the scalar field, which determines the expectation value of the boundary operator, can be solved analytically in this situation, and its scaling with \(\delta t\) can readily be seen from this solution. From this, we also obtain the scaling of the energy density in the boundary.

Consider a generic deformation of a conformal field theory (CFT) in \(d\) spacetime dimensions by the time-dependent coupling \(\lambda = \lambda(t)\) of a relevant operator \(\mathcal{O}_\Delta\) of dimension \(\Delta\): \(\mathcal{L}_0 \to \mathcal{L} = \mathcal{L}_0 + \lambda \mathcal{O}_\Delta\). The gravity
dual describing such a deformation is given by

\[ I_{d+1} = \frac{1}{16\pi G_{d+1}} \int d^{d+1}x \sqrt{-g} \times \left( R + d(d-1) - \frac{1}{2} (\partial \phi)^2 - \frac{1}{2} m^2 \phi^2 - u(\phi) \right), \tag{1} \]

where we have chosen an AdS radius of 1. The bulk scalar \( \phi \) is dual to \( \mathcal{O}_\Delta \) with \( m^2 = \Delta(\Delta - d) \). The potential \( u(\phi) \) contains terms of order \( \phi^3 \) or higher. To simplify our discussion, we will consider quenches where the conformal dimension of the operator is non-integer (for even \( d \) and not half-integer for odd \( d \) — see comments below). Further, we initially consider dimensions in the range \( \frac{d}{2} \leq \Delta < d \).

Since we are interested in quenches that are homogeneous and isotropic in the spatial boundary directions, we assume that both the background metric and the scalar field depends only on a radial coordinate \( \rho \) and a time \( t \). We will work in a spacetime asymptotic to the AdS Poincaré patch as \( \rho \to 0 \). Hence the bulk metric is

\[ ds^2 = -A(t, \rho)dt^2 + \Sigma(t, \rho)^2 d\tilde{\Omega}^2 + \rho^{-4} A(t, \rho)^{-1} d\rho^2. \tag{2} \]

The (nonlinear) Einstein equations and the scalar field equation then take the form:

\[
0 = -\frac{2(d-3)}{(d-1)A} u(\phi) + \frac{2d(d-3)}{A} \rho^4 (\phi')^2 - \frac{d-3}{(d-1)A} m^2 \phi^2 - \left( \frac{\dot{\phi}}{A} \right)^2 + 2(d-2)(d-1) \left[ \left( \frac{\dot{\Sigma}}{A \Sigma} \right)^2 - \left( \frac{\rho^2 \Sigma'}{\Sigma} \right)^2 \right],
\]

\[
0 = \frac{d}{(d-1)} - \frac{m^2 \phi^2}{2(d-1)} + \frac{\rho^4 A}{2(d-1)} (\phi')^2 + \frac{\ddot{\phi}}{A} - \rho^4 A' \Sigma' \rho^2 - \rho^4 A' \Sigma' + \rho^4 A^2 \Sigma, \tag{3}
\]

\[
0 = -\rho^2 \left( \frac{(\phi')^2}{2(d-1)} + \frac{1}{2(d-1)} \left( \frac{\dot{\phi}}{\rho^2 A} \right)^2 + \frac{\Sigma'}{\rho} + 2 \frac{\Sigma'}{\rho} \right), \tag{4}
\]

\[
0 = \rho^2 \frac{\phi' \dot{\phi}}{d-1} + \frac{\dot{A} \Sigma'}{A \Sigma} - \frac{A' \Sigma}{A \Sigma} + 2 \frac{\Sigma'}{\Sigma}, \tag{5}
\]

\[
0 = -\frac{\delta u(\phi)}{\delta \phi} - m^2 \phi + \rho^4 A \phi'' + 2 \rho^3 A \phi' + \rho^4 A' \phi' + \frac{(d-1)\rho^4 A \phi'}{\Sigma} + \frac{\dot{A} \phi}{A^2} - \frac{(d-1) \Sigma' \phi}{\Sigma} - \frac{\dot{\phi}}{A}. \tag{6}
\]

Here \( C = a_{d-2}(\infty) \) is an integration constant. With \( d = 4 \), this expression matches that found in [20], using Eddington-Finkelstein coordinates.

In our quenches, the coupling to \( \mathcal{O}_\Delta \) is made time-dependent with a characteristic time \( \delta t \) as

\[ \lambda = \lambda \left( t/\delta t \right). \tag{7} \]

For general \( \delta t \), the response \( p_{2\Delta-d} \) in eq. cannot be solved analytically. However, as described in [18, 20], for large \( \delta t \) (adiabatic quenches), we can find a series solution for \( \phi \) in inverse powers of \( \delta t \) and in principle, we can solve for \( p_{2\Delta-d} \) analytically.

We now present a new analytic approach for the opposite limit of fast quenches. That is, for quenches where \( \delta t \) is much smaller than any other scale. As described above, to answer the question of how much work is done by the quench, we need only consider the interval \( 0 \leq t \leq \delta t \). Intuitively, we may expect that when \( \delta t \) is very short, there is no time for nonlinearity in the bulk equations to become important, i.e., for the metric to backreact on
the scalar.

To make this intuition manifest, we rescale the coordinates and fields by the parameter $\delta t$ considering their (leading) dimension in units of the AdS radius: $\tilde{\rho} = \delta t \, \rho$, $t = \delta t \, \tilde{t}$, $A = \tilde{A}/\delta t^2$, $\Sigma = \bar{\Sigma}/\delta t$ and $\phi = \delta t^d - \Delta \, \hat{\phi}$. With this rescaling, the limit $\delta t \to 0$ then removes the scalar from the Einstein equations \(5\) \(8\), while leaving the form of the Klein-Gordon equation \(7\) unchanged.

The coefficient $a_{d-2}$ controls the next-to-leading order term in $A$ at small $\rho$. As we will show, this coefficient scales as $\delta t^{d-2\Delta}$. Further in eq. \(9\), this coefficient is accompanied by a factor of $\rho^2$ and hence this term has an overall scaling of $\delta t^{2(d-\Delta)}$. Hence as long as we are considering a relevant operator, this term vanishes in the limit $\delta t \to 0$. The same is true of the subleading contributions in the expression of $\bar{\Sigma}$. Hence for fast quenches with small $\delta t$, we can approximate the metric coefficients as simply

$$\bar{\Sigma} = \tilde{\rho}^{-1}, \quad \bar{A} = \tilde{\rho}^{-2}. \quad (12)$$

The equation for $\hat{\phi}$ becomes the Klein-Gordon equation in the AdS vacuum spacetime, i.e.,

$$\tilde{\rho}^2 \tilde{\rho}^2 \tilde{\rho}^2 \hat{\phi} - (d - 1) \tilde{\rho} \partial_t \hat{\phi} - \rho^2 \tilde{\rho}^2 \hat{\phi} + \Delta (d - \Delta) \hat{\phi} = 0. \quad (13)$$

That is, in the limit of small $\delta t$, the work done in the full nonlinear quench can be determined by simply solving the linear scalar field equation \(13\) in empty AdS space!

Now we consider sources that vanish for $t < 0$ and are constant for $t \geq \delta t$. In $0 < t < \delta t$, we vary the source as

$$p_0(t) = \delta p \left(\frac{t}{\delta t}\right)^\kappa \quad (14)$$

where $\kappa$ is a positive exponent. Note that here $p_0(t) \geq \delta t = \delta p$. Since $\phi = 0$ before we switch on the source at $t = 0$, it remains zero throughout the bulk up to the null ray $\tilde{t} = \rho$. Therefore we impose

$$\phi(t = \rho, \rho) = 0. \quad (15)$$

Evaluating the scalar field equation \(13\) subject to the boundary conditions \(14\) and \(15\), we find \(21\)

$$p_{2\Delta-d}(t) = b_\kappa \, \delta t^{d-2\Delta} \, \delta p \left(\frac{t}{\delta t}\right)^{d-2\Delta + \kappa} \quad (16)$$

with

$$b_\kappa = -\frac{2^{d-2\Delta} \, \Gamma(\kappa + 1) \, \Gamma\left(\frac{d+2}{2} - \Delta\right)}{\Gamma(d+1+\kappa-2\Delta) \, \Gamma\left(\frac{d+1}{2} - d\right) \, \Gamma\left(\frac{d+3}{2} - \Delta\right)}. \quad (17)$$

Of course, if we construct more complicated sources with a series expansion of monomials as in eq. \(13\), then since eq. \(13\) is linear, the response is simply given by the sum of corresponding terms as in eq. \(16\). As an example, consider the source

$$p_0(t) = 16 \, \delta p \left(\tilde{t}^2 - 2 \tilde{t}^3 + \tilde{t}^4\right) \quad (18)$$

as shown in fig. 2. In this case, the source vanishes in both the initial and final state and it reaches the maximum

![FIG. 2: Normalized source $p_0/\delta p$ for eq. \(18\) as a function of the rescaled time $\tilde{t} = t/\delta t$.](image)

$\delta p$ at $t = \delta t/2$. Figs. 3 and 4 show the corresponding response for various values of $\Delta$ in $d = 4$.

The response coefficient \(19\) exhibits two noteworthy features: First, we see that the overall scaling of the response is $\delta t^{d-2\Delta}$. This is precisely the behaviour found in the numerical studies of \(20\) in the case $d = 4$. Second of all, $p_{2\Delta-d}$ varies in time as $t^{d+\kappa-2\Delta}$. Therefore if $\kappa < 2\Delta - d$, the response (i.e., the operator expectation value $\langle O \rangle$ in the boundary theory) diverges at $t = 0$! For a source constructed as a series, both of these features in the response are controlled by the smallest exponent, as illustrated in figs. 3 and 4 for eq. \(18\).

For homogeneous quenches, the diffeomorphism Ward identity reduces to $\partial_t E = -\langle O \rangle \partial_t \lambda \quad 18 \quad 20$. Hence we can evaluate change in the energy density as

$$\Delta E = -\mathcal{A} E \int_{-\infty}^{+\infty} p_{2\Delta-d} \partial_t p_0 \, dt, \quad (19)$$

with \(22\)

$$\mathcal{A} E = \frac{2\Delta - d}{16\pi G_{d+1}} = \frac{(2\Delta - d)\pi^{d/2} \Gamma(\Delta/2)}{2d(d+1)\Gamma(d-1)} \, C_T. \quad (20)$$

Since $\partial_t p_0$ vanishes for $t < 0$ and $t > \delta t$, the above integral reduces to an integral from 0 to $\delta t$. It is for this reason that we do not need to determine the response $p_{2\Delta-d}$ after $t = \delta t$. Further, for fast quenches, the change in energy density will scale as $\delta t^{d-2\Delta}$. Note that $\partial_t p_0$ scales as $\delta t^{-1}$, but the range of the integral $0 < t < \delta t$ adds an additional scaling of $\delta t^{+1}$. Hence the net scaling of
$\Delta \mathcal{E}$ is precisely the scaling of $p_{2\Delta-d}$. Again this precisely matches the scaling found numerically in \cite{20} for $d = 4$. In fact, this behavior can be fixed as follows: Since eq. \cite{13} is linear, we must have $p_{2\Delta-d} \propto \delta \rho$ and hence $\Delta \mathcal{E} \propto \delta \rho^3$ from eq. \cite{19}. Finally, dimensional analysis demands $\Delta \mathcal{E} \propto \delta \rho^2/\delta t^{2\Delta-d}$, up to numerical factors.

However, recall the singular behaviour in the response at $t = 0$ for $\kappa < 2\Delta - d$. Despite this divergence, one can easily see that in fact, the corresponding integral \cite{19} remains finite as long as $\kappa > \Delta - \frac{d}{2}$. That is, for fixed $\Delta$ and $d$, we are constrained as to how quickly the source may be turned on. In fact, a more careful examination \cite{21} of the bulk solutions indicates that our analysis is valid for $\kappa > \Delta - \frac{d}{2} + \frac{1}{2}$. For quenches not satisfying this inequality, we can no longer ignore the backreaction of the scalar on the spacetime geometry.

To summarize, we have showed that in the limit of fast, abrupt quenches, the response and the energy density of a strongly coupled system which admits a dual gravitational description scales as $\delta t^{d-2\Delta}$. Here $\frac{d}{2} \leq \Delta < d$ is the conformal dimension of the quenched operator in the vicinity of the ultraviolet fixed point. Although we considered a quench from a vacuum state at $t = 0$, our results are universal. That is, they are independent of the initial state of the system, e.g., we may start with a thermal state, as in \cite{13,20}. This is again a reflection of the fact that abrupt holographic quenches are completely determined by the UV dynamics of the theory — see fig. 4. Also, if different operators are quenched simultaneously, the response is dominated by the one with the largest conformal dimension.

We emphasize that while our calculations only considered the linearized scalar equation \cite{13}, our results apply for the full nonlinear quench. In the limit $\delta t \rightarrow 0$, the relevant physics occurs in the far asymptotic geometry (see fig. 11) where the bulk scalar and perturbations of the AdS metric are all small. This contrasts with \cite{18,20}, which only worked within a perturbative expansion in the amplitude of the scalar. Of course, the scalings determined there match those found here, but it was uncertain if they would persist in a full nonlinear analysis.

Of course, the present analysis does not predict the dynamical evolution of the system for $t > \delta t$, however, we can deduce the equilibrium thermal state of the system as $t \rightarrow \infty$. Indeed, since the coupling and energy density are constant for $t > \delta t$, $\lambda(\pm \infty) = \lambda(\delta t)$ while eq. \cite{19} determines the final energy density of the system, to leading order in $\delta t$. Together, these parameters completely specify the final equilibrium state.

Note that our analysis strictly applies to relevant operators, for which $d - \Delta > 0$. With a marginal operator (i.e., $\Delta = d$), we can expect $\Delta \mathcal{E} \propto \delta t^{d-\Delta}$ on purely dimensional grounds \cite{16}. While this matches the scaling found above, our numerical coefficients would no longer be valid. Marginal operators were also considered in \cite{9,14} with a four-dimensional bulk. This case is analytically accessible because the scalar propagates on the light-cone. Extending this analysis to an odd-dimensional bulk is more challenging \cite{9} because the scalar propagator is nonvanishing throughout the interior of the light-cone, similar to that for the relevant operators studied here.

Our discussion was also limited to $\frac{d}{2} \leq \Delta < d$, while unitarity bounds also allow for $\frac{d}{2} - 1 \leq \Delta < \frac{d}{2}$. In the latter range, we must consider the so-called ‘alternate quantization’ of the bulk scalar \cite{21}. But, the asymptotic expansion of the scalar takes precisely the same form as in eq. \cite{5}. However, in this regime, $p_0 (p_{2\Delta-d})$ is the coefficient of the (non-)normalizable mode. Our analysis applies equally well for this range of $\Delta$ and so one still finds $p_{2\Delta-d} \propto \delta \rho \delta t^{d-2\Delta}$. That is, the response becomes vanishingly small as $\delta t \rightarrow 0$ with $\delta \rho$ kept fixed. Hence to produce a finite $(\mathcal{O}_\Delta)$ or finite $\Delta \mathcal{E}$, we would need to scale $\delta \rho$ with an inverse power of $\delta t$.

When $\Delta$ is an integer for even $d$ or half-integer for odd $d$, the scaling of the response $(\mathcal{O}_\Delta)$ receives additional log($\delta t$) corrections \cite{18}. These logarithmic corrections arise from log $\rho$ modifications in the asymptotic expansion \cite{5} of the bulk scalar and are easily computed analytically following the present approach \cite{21}.

Another exceptional case arises with $\kappa = 2\Delta - d - n$ where $n$ is a positive integer. In this case, eq. \cite{17} indicates $b_n = 0$. Hence if the source is given by a series of monomials \cite{14}, the scaling of the response will be controlled by the first subleading contribution. With a single monomial, the (subleading) scaling of the response is controlled by nonlinearities in the bulk equations \cite{21}, i.e., $p_{2\Delta-d} \sim \delta t^{-\Delta}(\delta \rho \delta t^{d-\Delta})^n$ where $n = 2$ if the potential contains a $\delta^3$ term and $n = 3$ otherwise.

It is interesting to consider the limit of abrupt quenches with $\delta t = 0$, as this usually sets the starting point in analyses at weakly coupling. Our holographic result, $\Delta \mathcal{E} \propto \delta \rho^2/\delta t^{2\Delta-d}$, indicates that the energy density diverges for an abrupt quench with $\Delta > \frac{d}{2}$ (a logarithmic divergence appears for $\Delta = \frac{d}{2}$ \cite{13,21}). Hence it would be interesting to carefully compare these holographic results with those for the weak coupling calculations of, e.g., \cite{2,6,24}. Let us note here that certain singular behaviours were observed for abrupt quenches of a fermionic mass term \cite{24}. Of course, the preceding considerations assume $\delta \rho$ is held fixed in the limit $\delta t \rightarrow 0$. Instead, if we scale the source to zero as $\delta \rho \propto \delta t^{d-\frac{d}{2}}$, $\Delta \mathcal{E}$ will remain
finite. However, we stress that this limit still produces a divergent response since $p_{2\Delta-d} \sim d \delta t^{d-2\Delta} \propto \delta t^{\frac{d}{2}-\Delta}$.

An important question to ask is to what extent our results are relevant for everyday physical systems. Gauge theories with a dual gravitation description are necessarily strongly coupled and have an ultraviolet fixed point with large central charge. The framework of the gauge-string duality allows for the study of both the finite 't Hooft coupling corrections (the higher-derivative corrections in the gravitational dual) and non-planar (quantum string-loop) corrections. We expect that our gravitational analysis are robust with respect to the former, as the relevant near-boundary space-time region is weakly curved. Whether finite central charge corrections are important or not is an open question.

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