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# Transient Orthogonality Catastrophe in a Time-Dependent Nonequilibrium Environment

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We study the response of a highly excited time-dependent quantum many-body state to a sudden local perturbation, a sort of orthogonality catastrophe problem in a transient nonequilibrium environment. To this extent we consider, as a key quantity, the overlap between time-dependent wave functions, which we write in terms of a novel two-time correlator generalizing the standard Loschmidt echo. We discuss its physical meaning, general properties, and its connection with experimentally measurable quantities probed through nonequilibrium Ramsey interferometry schemes. Then we present explicit calculations for a one-dimensional interacting Fermi system brought out of equilibrium by a sudden change of the interaction, and perturbed by the switching on of a local static potential. We show that different scattering processes give rise to remarkably different behaviors at long times, quite opposite from the equilibrium situation. In particular, while the forward scattering contribution retains its power-law structure even in the presence of a large nonequilibrium perturbation, with an exponent that is strongly affected by the transient nature of the bath, the backscattering term is a source of nonlinearity which generates an exponential decay in time of the Loschmidt Echo, reminiscent of an effective thermal behavior.

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**Introduction.**—The response of gapless quantum many-body systems to sudden local perturbations is a remarkably nonlinear phenomenon; even a weak disturbance substantially changes the structure of the many-body state. Signatures of this *orthogonality catastrophe* (OC) emerge in various condensed matter settings [1], from x-ray spectra in metals [2] and Luttinger liquids (LLs) [3–6] to the physics of the Kondo effect [7–11], and typically results in power-law decays of dynamical correlations. Recently, impressive experimental developments with ultracold atomic gases [12] have made it possible to create and probe local excitations in a quantum many-body system with single-site and real-time resolution [13,14], bringing fresh new input to this venerable problem [15]. While most of the attention has been traditionally devoted to perturbations acting on systems in their ground state or—more recently—in driven stationary nonequilibrium conditions [16–22], much less is known about the response of explicitly time-dependent quantum states, such as, for example, those obtained by rapidly changing in time some parameter of an otherwise isolated system. The problem is of current experimental relevance since ultracold gases have proven to be natural laboratories where dynamical quantum correlations can be probed in the time domain. In addition, it also raises a number of intriguing theoretical questions. A coherent time-dependent excitation, such as a sudden global quench, creates an effective nonequilibrium time-dependent bath for the local degrees of freedom. What is the effect of such an environment on the OC phenomenon and its associated power laws? For a generic, nonintegrable,

quantum many-body system one might expect this environment to be, at sufficiently long times, effectively thermal, turning the power-law decay of the OC correlator into an exponential. Yet, strongly interacting quantum systems may often get trapped into long-lived metastable prethermal states which may still show genuine quantum correlations [23–27]. Investigating the local spectral properties of these transient states of nonequilibrium quantum matter is among the purposes of this Letter.

**Transient OC protocol.**—We begin with a general discussion of the nonequilibrium protocol that will be the focus of this Letter. We consider a quantum many-body system initially prepared at time  $t_0 = 0$  in the ground state  $|\Psi_0\rangle$  of some Hamiltonian  $\mathcal{H}_0$ . We then let the system evolve up to some time  $t > t_0$  with a new Hamiltonian  $\mathcal{H}$ , i.e.,  $|\Psi(t)\rangle = e^{-i\mathcal{H}t}|\Psi_0\rangle$ , which differs from  $\mathcal{H}_0$  by a sudden change of some global parameter. This global quantum quench injects extensive energy into the system and triggers transient nonequilibrium dynamics. To gain insights into the structure of this transient state we will then switch on a local perturbation  $V_{\text{loc}}$  for an interval  $\tau$  between  $t$  and  $t' = t + \tau$  and compare the states obtained at time  $t'$ , respectively, in the presence or absence of the local perturbation. In the spirit of an OC problem we study the overlap

$$\mathcal{D}(t', t) \equiv \langle \Psi(t') | \Psi_{t+}(t') \rangle, \quad (1)$$

where  $|\Psi_{t+}(t')\rangle \equiv e^{-i\mathcal{H}_+(t'-t)}|\Psi(t)\rangle$ , with  $\mathcal{H}_+ = \mathcal{H} + V_{\text{loc}}$ . This overlap can be written suggestively as a two-time dynamical correlator,

$$\mathcal{D}(t', t) = \langle \Psi(t) | e^{i\mathcal{H}(t'-t)} e^{-i\mathcal{H}_+(t'-t)} | \Psi(t) \rangle. \quad (2)$$

One immediately sees that when the initial state  $|\Psi_0\rangle$  is the ground state of  $\mathcal{H}$ ; then,  $\mathcal{D}(t', t) \equiv D_{\text{eq}}(\tau) = \langle \Psi_0 | e^{i\mathcal{H}\tau} e^{-i\mathcal{H}_+\tau} | \Psi_0 \rangle$ ; i.e., it becomes time-translational invariant and reduces to the familiar OC correlator, also known as the core-hole Green's function in the x-ray edge problem [2] or the Loschmidt echo amplitude [28–31], which recently has attracted renewed interest in the context of work statistics [32–36]. In equilibrium, the large-time asymptotics of  $D_{\text{eq}}(\tau)$  gives rich information on the ground state  $|\Psi_0\rangle$  and its low-lying excitations. In particular a power-law decay reflects an orthogonality catastrophe in the low-energy sector induced by the local perturbation, mirroring the one introduced by Anderson for stationary (ground) states [1]. Considering the enormous principle importance of this phenomenon for equilibrium quantum many-body physics, it is natural to investigate its fate for time-dependent excited states, as we are going to do in the following. Before turning to an example, it is useful to discuss some general features of our transient OC correlator. We start by writing it in terms of the exact eigenstates of  $\mathcal{H}$ ,  $\mathcal{H}|\Phi_n\rangle = E_n|\Phi_n\rangle$  [37],

$$D(\tau; t) = \sum_{nm} \rho_{nm}(t) \langle \Phi_m | e^{i\mathcal{H}\tau} e^{-i\mathcal{H}_+\tau} | \Phi_n \rangle, \quad (3)$$

where  $\rho_{nm}(t) = \langle \Phi_n | \Psi(t) \rangle \langle \Psi(t) | \Phi_m \rangle$ . Differently from the equilibrium case, here both diagonal and off-diagonal matrix elements contribute to  $D(\tau; t)$  with a time-dependent amplitude  $\rho_{nm}(t)$  encoding information about the state of the system before the switching on of the local perturbation. The result highlights the very nature of our transient correlator as a sensitive probe of OC in excited many-body states. By averaging Eq. (3) over the waiting time  $t$  and taking the Fourier transform with respect to  $\tau$  we obtain [37]  $P(W) \equiv \int d\tau e^{iW\tau} D(\tau; t)$  as

$$P(W) = \sum_{n\alpha} \delta(W - \tilde{E}_\alpha + E_n) |\langle \Phi_n | \tilde{\Phi}_\alpha \rangle|^2 \rho_{nn}(0), \quad (4)$$

where we have introduced the eigenstates of the full Hamiltonian,  $\mathcal{H}_+|\tilde{\Phi}_\alpha\rangle = \tilde{E}_\alpha|\tilde{\Phi}_\alpha\rangle$ , in the presence of the local perturbation.  $P(W)$  is the probability distribution of the work done, starting from a *nonequilibrium state* described by the diagonal ensemble  $\rho_{nn}$ , and suddenly switching on the local potential. This extends to the non-equilibrium case the connection between the OC correlator and work statistics and represents an interesting result on its own, in view of recent theoretical interest in characterizing the work statistics and its universal properties in nonthermal ensembles [38,39]. While it is known that in equilibrium the work statistics of local quenches shows zero-temperature edge singularities [9,32], we will see below how this is modified in the presence of bulk excitations.

Finally, it is interesting to discuss experimental protocols to measure the transient OC correlator. Recent proposals

outline how the equilibrium OC correlator may be measured [15,31,40–45]. The key is to use an auxiliary two-level system (TLS) or qubit, coupled to the system through the local perturbation,  $\mathcal{H}[\sigma^z] = \mathcal{H} + V_{\text{loc}}(1 + \sigma^z)/2$ . Cold atoms and other hybrid systems, such as circuit QED units, represent the natural platforms to realize this. Here we extend these ideas to design nonequilibrium Ramsey interferometry schemes to manipulate the TLS in such a way as to obtain  $D(t'; t)$  out of simple local TLS measurements. For example, the real part of the transient OC correlator can be obtained by measuring the  $x$  component of TLS magnetization [37],

$$\text{Re}D(t'; t) = \langle \Phi_t(t') | \sigma^x | \Phi_t(t') \rangle, \quad (5)$$

where  $|\Phi_t(t')\rangle$  is a state obtained with a specific protocol involving (i) a first evolution up to time  $t$ , (ii) a Ramsey  $\pi/2$  pulse, and (iii) a second time evolution up to time  $t' = t + \tau$ . Alternatively, one can obtain  $D(t'; t)$  from a TLS dynamical correlator, a nonequilibrium analog of the core-hole Green's function in the x-ray edge problem [37].

*Quenched Luttinger liquid in a transient local potential.*—We now focus our attention on a one-dimensional spinless interacting Fermi system described by the Luttinger model [46] and brought out of equilibrium by a sudden quench of the interaction. In the bosonization language, the Hamiltonian of the system after the quench can be written in terms of collective LL degrees of freedom  $\phi(x)$ ,  $\theta(x)$  as

$$\mathcal{H} = \frac{u}{2\pi} \int dx \left[ K(\partial_x \theta(x))^2 + \frac{1}{K}(\partial_x \phi(x))^2 \right], \quad (6)$$

where the initial state  $|\Psi_0\rangle$  corresponds to the ground state of Eq. (6) with the Luttinger parameter  $K_0 \neq K$  [62]. The dynamics after this global quantum quench has been studied in great detail [47]. Since we are interested in the correlator defined in Eq. (2) we have to discuss the nature of the local perturbation. Here we consider a static potential that couples to the electron density through a forward (fs) and a backward (bs) scattering term [46,48] which in bosonic variables is

$$V_{\text{loc}} \equiv V_{\text{fs}} + V_{\text{bs}} = g_{\text{fs}} \partial_x \phi(x)|_{x=0} + g_{\text{bs}} \cos 2\phi(x=0). \quad (7)$$

The evaluation of the transient OC correlator  $\mathcal{D}(t', t)$  greatly simplifies by noticing that, as in equilibrium, the forward- and backward-scattering processes are decoupled; i.e., they involve independent degrees of freedom [37]. As a result, we find that the OC correlator factorizes into  $\mathcal{D}(t', t) = \mathcal{D}_{\text{fs}}(t', t) \mathcal{D}_{\text{bs}}(t', t)$ .

*Forward scattering.*—Let us start by discussing the forward-scattering contribution, which can be computed exactly using a method due to Schotte and Schotte [48,49]. The final result is

$$\mathcal{D}_{\text{fs}}(t', t) = \langle T e^{-i\eta\theta(0,t')} e^{i\eta\theta(0,t)} \rangle, \quad (8)$$

with  $\eta = g_{\text{fs}}K/u$ , which can be evaluated in terms of local correlators of the quenched LL [37]. It is useful to write the OC correlator as a function of the variables  $t' - t = \tau$  and  $t$ ,

$$D_{\text{fs}}(\tau; t) = \frac{e^{-i\varphi(\tau)}}{[1 + (\Lambda\tau)^2]^{\delta_{\text{neq}}^{\text{oc}}/2}} f_t(\tau), \quad (9)$$

where  $\Lambda$  is an ultraviolet cutoff and the transient function  $f_t(\tau)$  is

$$f_t(\tau) = \left( \frac{[1 + \Lambda^2(2t + \tau)^2]^2}{[1 + (2\Lambda t)^2][1 + 4\Lambda^2(t + \tau)^2]} \right)^{\delta_{\text{tr}}^{\text{oc}}/4}, \quad (10)$$

where the phase  $\varphi(\tau) = \frac{g_{\text{fs}}^2}{2u^2} K \text{sgn}(\tau) \arctan(\Lambda\tau)$ , while the two exponents read, respectively,  $\delta_{\text{neq/tr}}^{\text{oc}} = (g_{\text{fs}}^2/2u^2) K_{\text{neq/tr}}$  with  $K_{\text{neq/tr}} = K_0(1 \pm K^2/K_0^2)/2$ . We immediately see that for  $K = K_0$  the above expression reduces to the well-known result for the equilibrium LL,  $D_{\text{fs}}^{\text{eq}}(\tau) \sim (1/\Lambda\tau)^{\delta_{\text{oc}}^{\text{eq}}}$ , with an OC exponent  $\delta_{\text{oc}}^{\text{eq}} = K_0 g_{\text{fs}}^2/2u^2$  [3,48,50]. In the case of a bulk quench,  $K_0 \neq K$  the situation is more interesting. As we see from Fig. 1 when both time arguments are longer than the microscopic time scale, i.e.,  $t, \tau \gg 1/\Lambda$ , the OC correlator features two distinct power-law regimes, with a crossover scale set by the total time  $t$  after the bulk quench. The intermediate-time regime,  $1/\Lambda \ll \tau \ll t$ , when the duration of the local perturbation is

much shorter than that of the global bulk excitation, can be described by a completely dephased nonequilibrium environment which gives a decay  $D_{\text{fs}} \sim \tau^{-\delta_{\text{neq}}^{\text{oc}}}$ . While this would be the leading power-law behavior for a strictly infinite waiting time, the transient nature of the environment results in a different power-law decay at longer times; i.e., for  $\tau \gg t$  we have  $D_{\text{fs}} \sim \alpha_t \tau^{-\delta_{\text{new}}^{\text{oc}}}$ , with a prefactor  $\alpha_t \sim 1/t^{\delta_{\text{tr}}^{\text{oc}}/2}$  that depends on the waiting time  $t$ . Here  $\delta_{\text{new}}^{\text{oc}} = \delta_{\text{neq}}^{\text{oc}} - \delta_{\text{tr}}^{\text{oc}}/2 = (K_{\text{neq}} - K_{\text{tr}}/2)g_{\text{fs}}^2/2u^2$  can be larger or smaller than the short-time OC exponent  $\delta_{\text{neq}}^{\text{oc}}$  depending on the sign of  $K_{\text{tr}}$ , i.e., whether  $K \gtrless K_0$ .

**Back scattering.**—We now consider the backscattering potential and start with a perturbative calculation of  $\mathcal{D}_{\text{bs}}(t', t)$ . Using the linked-cluster theorem we may write

$$\mathcal{C}(t', t) \equiv \log \mathcal{D}_{\text{bs}}(t', t) = \langle \Psi_0 | T e^{-i \int_t^{t'} dt_1 V_{\text{bs}}(t_1)} | \Psi_0 \rangle_c, \quad (11)$$

where only connected (c) averages contribute,  $V_{\text{bs}}$  is defined in Eq. (7), and  $V_{\text{bs}}(t)$  is in the interaction representation of the post-quench global Hamiltonian. Expanding in powers of  $V_{\text{bs}}$  we get to lowest order

$$\mathcal{C}^{(2)}(t', t) = -\frac{g_{\text{bs}}^2}{2} \int_t^{t'} dt_1 dt_2 \langle T [\cos(2\phi(t_1)) \cos(2\phi(t_2))] \rangle_c. \quad (12)$$

The result can be evaluated in terms of local correlators of the quenched LL [37]. In equilibrium,  $K_0 = K$ , the OC correlator is a function only of  $\tau = t' - t$  and the integral (12) can be evaluated analytically [48]. The result shows that for  $K > 1$  the correlator goes to a constant at long times; i.e., perturbation theory in  $V_{\text{bs}}$  is well behaved. In contrast, for  $K < 1$  the perturbative correction blows up at long times,  $\Lambda\tau \gg 1$ , as  $\text{Re}\mathcal{C}^{(2)}(\tau) \sim -(g_{\text{bs}}^2/\Lambda^2)(\Lambda\tau)^{2(1-K)}$ , consistent with the result that the backscattering potential is relevant in this regime [5]. A crossover time scale  $\tau_*^{\text{eq}}(K)$  can be extracted by setting  $\mathcal{C}^{(2)}(\tau_*^{\text{eq}}) \sim 1$ , to give  $\Lambda\tau_*^{\text{eq}} \sim (\Lambda/g_{\text{bs}})^{1/(1-K)}$ . In the strong-coupling phase ( $K < 1$ ) one can show that at long times ( $\tau \gg \tau_*^{\text{eq}}$ ) the correlator  $\mathcal{C}^{(2)}(\tau)$  can be resummed into a logarithmic divergence, which gives rise to a universal power law for  $\mathcal{D}_{\text{bs}}(\tau) \sim 1/\tau^{1/8}$  [3,51–56]. Finite-temperature effects generally change this power law to an exponential decay of the OC correlator [37].

Let us now study the nonequilibrium case,  $K_0 \neq K$ , which is considerably richer. We use the same parametrization as in the forward-scattering case, in terms of  $\tau = t' - t$  and the total time  $t$  after the quench, i.e.,  $\mathcal{C}_t^{(2)}(\tau)$ . The integral (12) cannot be evaluated in closed analytical form for  $K_0 \neq K$ , yet the structure of the solution can be understood from numerics. To simplify the discussion let us assume first an infinite waiting time after the quench, i.e.,  $t \rightarrow \infty$ , at fixed  $\tau$ . In Fig. 1 we plot  $-\text{Re}\mathcal{C}_\infty^{(2)}(\tau)$ , which reveals two different behaviors depending on the values of

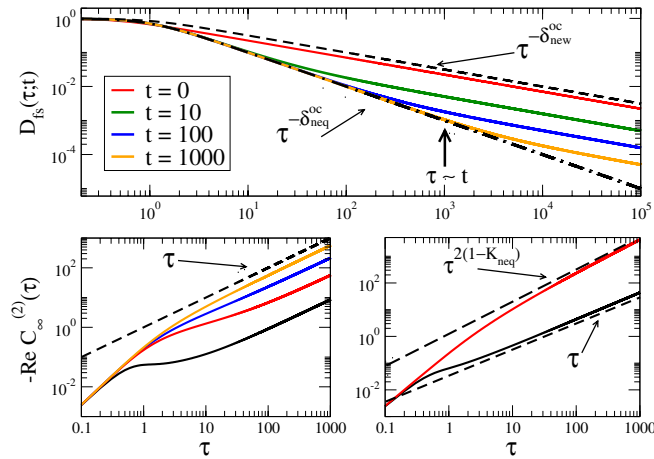


FIG. 1 (color online). Top panel: Forward-scattering contribution to the OC correlator  $D_{\text{fs}}(\tau; t)$ , for  $K_0 = 4$ ,  $K = 0.1$ . The dynamics features two distinct power-law regimes: a short-time regime with exponent  $\delta_{\text{neq}}^{\text{oc}}$  and a long-time one with exponent  $\delta_{\text{new}}^{\text{oc}}$ , with the crossover controlled by the total time  $t$ . Bottom panel: Lowest-order back-scattering contribution to the OC correlator,  $-\text{Re}\mathcal{C}_\infty^{(2)}(\tau)$ . In the left panel for  $K_0 = 1.5$  and  $K = 0.5, 0.75, 1.0, 2.5$  (top to bottom) corresponds to the *thermal* regime, with a linear power-law divergence. In the right panel the transition to the strong coupling regime for  $K_0 = 0.75$  and  $K = 0.25, 2.5$  (top to bottom), respectively, below and above  $K_* \approx 0.43$ .

$K_0$  and  $K$ . For  $K_0 > 1$  or  $K_0 < 1$  and  $K > K_* = \sqrt{K_0(1-K_0)}$  (see lower right panel in Fig. 2), corresponding to  $K_{\text{neq}} > 1/2$ , the leading long-time behavior has a linear divergence  $\text{Re}C_\infty^{(2)}(\tau) \sim -\gamma_*\tau$  with a prefactor that can be evaluated in closed form, (all energy scales in units of  $\Lambda$ )

$$\gamma_* = g_{\text{bs}}^2 \left( \frac{2\pi}{2^{K_{\text{neq}}} [2K_{\text{neq}} - 1]} \right) \frac{\Gamma(2K_{\text{neq}})}{\Gamma(K_{\text{neq}} + K) \Gamma(K_{\text{neq}} - K)}, \quad (13)$$

where  $\Gamma(x)$  is the Gamma function. Such behavior results, through Eq. (11), in an exponential decay of the OC correlator  $\mathcal{D}_{\text{bs}}(\tau) \sim \exp(-\gamma_*\tau)$ , implying a Lorentzian work distribution  $P(W)$  of width  $\gamma_*$ , which is reminiscent of an equilibrium finite-temperature behavior [37].

To further investigate this regime we define a time-dependent relaxation rate  $\gamma(\tau) = -\partial_\tau \text{Re}C_\infty^{(2)}(\tau)$  and plot it in Fig. 2 (top panels) for different values of  $K_0$ ,  $K$ . From this we see that upon approaching  $K \rightarrow K_*$  (right panel) the relaxation rate diverges and—quite differently from the thermal equilibrium case—we find a region of parameters for  $K_0 < 1$  and  $K < K_*$  (see lower right panel in Fig. 2), corresponding to  $K_{\text{neq}} < 1/2$ , where the correlator  $C_\infty^{(2)}(\tau)$  still diverges as a power law at long times but with a new exponent,  $\text{Re}C_\infty^{(2)}(\tau) \sim -(g_{\text{bs}}^2/\Lambda^2)(\Lambda\tau)^{2(1-K_{\text{neq}})}$ . Such a divergence suggests that the problem retains in this regime a strong-coupling nature and allows us to define a non-equilibrium crossover time scale  $\Lambda\tau_*^{\text{neq}} \sim (\Lambda/g_{\text{bs}})^{1/(1-K_{\text{neq}})}$  controlling the flow of the backscattering potential (see below). While these results have been obtained assuming an infinite waiting time after the quench, finite-time effects do not seem to qualitatively change this behavior [37].

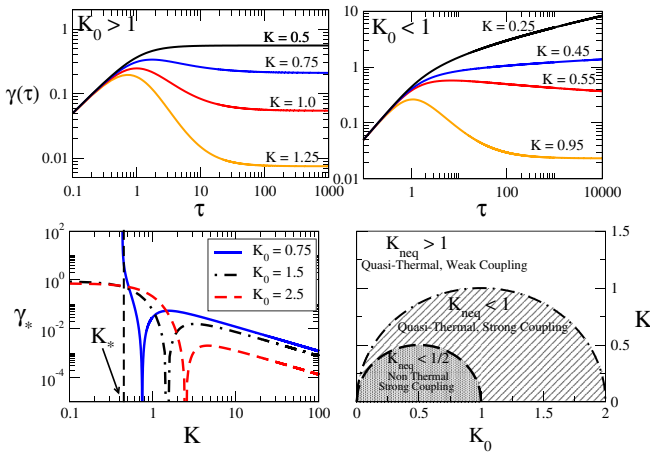


FIG. 2 (color online). Top panels: Time-dependent relaxation rate  $\gamma(\tau)$  for  $K_0 = 1.5$  (left panel) and  $K_0 = 0.75$  (right panel) and different values of  $K$ . For  $K_0 < 1$  and upon approaching  $K \rightarrow K_* = 0.43$  the relaxation rate diverges. Bottom panels: Long-time limit of the relaxation rate  $\gamma(\tau) \rightarrow \gamma_*$  as a function of  $K$ ,  $K_0$  with  $\gamma_* \sim g_{\text{bs}}^2 (K_0 - K)^2$  for  $K \rightarrow K_0$  (left panel), and a summary of different dynamical regimes (right panel).

We now consider the effect of higher-order backscattering terms using a time-dependent renormalization group (RG) method recently developed for the bulk Sine-Gordon problem [57]. We refer the reader to the Supplemental Material [37] for details about the derivation of the RG flow, and here we discuss the main results. At two-loop order we obtain renormalization corrections to the vertex  $g_{\text{bs}}$  as well as to the quadratic part of the action, which we parametrize in terms of a dissipation  $\eta$  (an effective friction for the local coordinate due to its coupling to the bulk degrees of freedom) and an effective temperature  $T_{\text{eff}}$  [37,58]. Setting  $\Lambda = 1$ , we obtain flow equations,

$$\frac{dg_{\text{bs}}}{d \ln l} = g_{\text{bs}} \left[ 1 - \left( K_{\text{neq}} + \frac{K_{\text{tr}}}{1 + 4T_m^2} \right) \right] \quad (14)$$

$$\frac{d\eta}{d \ln l} = 2g_{\text{bs}}^2 I_\eta(T_m), \quad (15)$$

$$\frac{d(\eta T_{\text{eff}})}{d \ln l} = \eta T_{\text{eff}} + g_{\text{bs}}^2 I_{T_{\text{eff}}}(T_m); \quad \frac{dT_m}{d \ln l} = -T_m, \quad (16)$$

where  $T_m$  is the time after the quench, and  $I_\eta(T_m)$ ,  $I_{T_{\text{eff}}}(T_m)$  are given in the Supplemental Material [37]. The initial conditions for the RG are  $\eta(l=1) = 2/(\pi K)$ ,  $T_{\text{eff}}(l=1) = 0$ . If we take the long-time ( $T_m \rightarrow \infty$ ) limit, from the RG flow we immediately see the following: (i) for  $K_{\text{neq}} > 1$  the backscattering is irrelevant, and yet it generates an effective temperature whose value is  $T_{\text{eff}} \sim g_{\text{bs}}^2 I_{T_{\text{eff}}}(T_m = \infty)/\eta(l=1)$ ; (ii) for  $K_{\text{neq}} < 1$  the problem flows to strong coupling on time scales larger than  $\tau_*^{\text{neq}}$ , and yet a well defined effective temperature can still be identified, at least as long as the relaxation rate  $\gamma_*$  or the local dissipation  $\eta$  stays finite, i.e., for  $K_{\text{neq}} > 1/2$  (see Fig. 2). Deep in the strong-coupling phase the perturbative analysis suggests that the OC correlator might eventually keep its power-law behavior; this result would be remarkable, but from the current analysis we cannot firmly conclude whether a genuine nonequilibrium strong-coupling regime remains intact or if higher-order corrections eventually cut the RG flow.

We note that the effective-temperature  $T_{\text{eff}}$  is not equivalent to a true temperature  $T$ , since the latter implies a relaxation rate in the OC correlator  $\gamma_T \propto g_{\text{bs}}^2 T^{2K-1}$  (see Supplemental Material [37]), while the former gives a relaxation rate  $\gamma_* \propto T_{\text{eff}}/(2K_{\text{neq}} - 1)$ . Thus even in the long-time limit, and for a local nonlinearity, important differences arise between the OC physics in the presence of a thermalized bath and the one studied in this paper, where the bath is in a nonequilibrium prethermalized state. It is interesting to note that an effective temperature is also generated in very different contexts, such as nonequilibrium steady states arising due to externally imposed voltage bias [20] and noise [22].

**Conclusions.**—We have introduced a nonequilibrium OC problem, highlighting its experimental relevance,

and have studied it in detail for one-dimensional interacting fermions out of equilibrium due to a global quench. We have discussed how the nonequilibrium environment affects the forward- and backward-scattering contribution to the OC correlator. We have shown that, while the former retains its power-law structure even in the presence of a strong nonequilibrium perturbation, the latter is a source of nonlinearity that generates locally an exponential decay in time of the OC correlator. Such a result appears to be consistent with a recent numerical study on Loschmidt echo decay in a one-dimensional spin chain after a local quench starting from a highly excited state [59].

We expect the physics of this *quench-induced* decoherence to be relevant in other contexts as well, most notably in steady-state transport, where it will result in a nonvanishing zero bias conductance in the weak-link limit as well as a nonzero backscattering correction in the dual limit [60], a result which is consistent with the nonvanishing impurity density of states recently found in Ref. [61].

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