Orthogonality Catastrophe in Dissipative Quantum Many-Body Systems

F. Tonielli,[1] R. Fazio,[2] S. Diehl,[1,3] and J. Marino[3,4]

[1]*Institut für Theoretische Physik, Universität zu Köln, D-50937 Cologne, Germany*

[2]*Abdus Salam ICTP, Strada Costiera 11, I-34151 Trieste, Italy and NEST, Scuola Normale Superiore and Istituto Nanoscienze-CNR, I-56126 Pisa, Italy*

[3]*Kavli Institute for Theoretical Physics, University of California, Santa Barbara, California 93106-4030, USA*

[4]*Department of Physics, Harvard University, Cambridge Massachusetts 02138, USA and Department of Quantum Matter Physics, University of Geneva, 1211 Geneve, Switzerland*

(Received 24 September 2018; revised manuscript received 9 January 2019; published 30 January 2019)

We present an analog of the phenomenon of orthogonality catastrophe in quantum many-body systems subject to a local dissipative impurity. We show that the fidelity $F(t)$, giving a measure for distance of the time-evolved state from the initial one, displays a universal scaling form $F(t) \propto t^{\theta} e^{-\gamma t}$, when the system supports long-range correlations, in a fashion reminiscent of traditional instances of orthogonality catastrophe in condensed matter. An exponential falloff at rate γ signals the onset of environmental decoherence, which is critically slowed down by the additional algebraic contribution to the fidelity. This picture is derived within a second-order cumulant expansion suited for Liouvillian dynamics, and substantiated for the one-dimensional transverse field quantum Ising model subject to a local dephasing jump operator, as well as for XY and XX quantum spin chains, and for the two-dimensional Bose gas deep in the superfluid phase with local particle heating. Our results hint that local sources of dissipation can be used to inspect real-time correlations and to induce a delay of decoherence in open quantum many-body systems.

DOI: 10.1103/PhysRevLett.122.040604

Introduction.—Anderson's orthogonality catastrophe (OC) [1] is a paradigm in solid state physics [2] highlighting the sensitivity of a gapless many-body ground state to static and dynamical local perturbations. An x-ray absorption process creates into an electron gas a core hole which acts as a static potential, provoking a catastrophic response in the system: the ground states of the electron gas, with and without the core-hole potential, are orthogonal—the overlap between the two scaling as a decaying power law of the system size. Singular features manifest in dynamical properties as well: the Green's function of the core hole has a power-law decay at long times, departing from a simple free-particle behavior; in frequency domain, close to the threshold energy, the x-ray absorption spectrum vanishes algebraically, signaling the suppression of absorption processes in this energy window [3]. Orthogonality catastrophe has been corroborated in a number of systems ranging from Luttinger liquids [4] to Kondo models [5,6] and disordered metals [7], and it has recently received novel attention [8–16], thanks to experimental progresses in cold gases, where local excitations can be created in a quantum many particle system at ease [17,18].

The connection between OC and the return probability, or Loschmidt echo [9,19–28] $\mathcal{L}(t)$, is a recent interesting development in this evergreen problem. The overlap between the unperturbed ground state of a quantum Ising chain at criticality $|\psi(0)\rangle$ and the same state evolving in the presence of a defect of strength δg along the transverse field direction $|\psi(t)\rangle$ exhibits an analogous algebraic scaling behavior [29] to the one discussed above, $\mathcal{L}(t) = |\langle\psi(0)|\psi(t)\rangle|^2 \propto t^{-\theta}$, with $\theta \propto (\delta g)^2$. The physical rationale behind the "catastrophe" stands in the underlying criticality of the many-body system upon which the perturbation is applied: the diverging characteristic correlation length and times at the critical point facilitate the spread of the local disturbance across the whole system, making possible the orthogonality among the initial state and the evolved one as time increases. This setup can also be extended to non-equilibrium closed environments [30]: the system is first sent out of equilibrium by a quantum quench of a global Hamiltonian parameter, and later subject to the action of a local potential, resulting in a two-times orthogonality catastrophe which may show such novel features as aging dynamics [31].

In this work, we demonstrate that the phenomenon of OC is not only exclusive to *unitary* dynamics, rather it can also occur in a gapless quantum many-body system when a local *noisy* or *dissipative* perturbation is suddenly switched; this dissipative analog of the OC is presented through a number of instances ranging from low-dimensional quantum spin chains to the Bose-Hubbard model in the superfluid phase. In particular, we show the emergence of a power law scaling in time for the fidelity (a proper analogue of the Loschmidt echo for generic mixed states) of a system with critical, or, in general, long-range correlations, in a fashion reminiscent of the OC in closed gapless systems.

However, contrary to traditional instances of OC, the additional algebraic contribution to the fidelity determines a critical slow down of decoherence. The paradigm shift presented here for the OC can be experimentally accessible as localized dissipations can be tailored in ultracold gases [32–39], with the long-run perspective to employ local dissipative channels to detect gapless modes in open quantum many-body systems.

Orthogonality catastrophe from a Lindbladian impurity.—To illustrate this concept, we consider as a minimal model the one-dimensional quantum Ising chain [40] in a transverse field, $H_0 = -(J/2)\sum_i[\hat{\sigma}_i^x\hat{\sigma}_{i+1}^x + g\hat{\sigma}_i^z]$, we prepare at time $t = 0$ the system in its critical ground state ($g = 1$), and suddenly switch at subsequent times $t > 0$ a spin-dephasing Lindblad operator, $\hat{L} = \hat{\sigma}_j^z$, acting on a given site j of the chain. Using a Jordan-Wigner transformation [40], the critical Ising chain can be mapped into a one-dimensional system of gapless, free fermions with a local dephasing noise, $\hat{L} \propto \hat{n}_j$, occurring at rate $\sqrt{\kappa}$, and proportional to the density n_j, of Jordan-Wigner fermions. The dynamics of this system is accordingly ruled by the quantum master equation,

$$\dot{\rho}(t) = -i[H_0, \rho(t)] + \kappa\mathcal{L}[\rho(t)], \tag{1}$$

where $\mathcal{L}[\rho(t)] = \hat{L}\rho(t)\hat{L}^\dagger - \frac{1}{2}\{\hat{L}^\dagger\hat{L}, \rho(t)\}$, and with $\hat{L} = \hat{L}^\dagger$, $\hat{L}^2 = 1$ in this specific case. The dynamics ruled by the quantum master equation with Hamiltonian $\hat{H}_0$ and with a single Hermitian Lindblad operator $\hat{L} = \hat{\sigma}_j^z$ is equivalent [41,42] to the stochastic Schrödinger evolution governed by

$$\hat{H}_\eta(t) = \hat{H}_0 + \sqrt{\kappa}\eta(t)\hat{L}, \tag{2}$$

where $\eta(t)$ is a Gaussian white noise and $\hat{L}$ is, for instance, a local spin perturbation along the transverse field direction, as in the case under study in this work. The Lindblad evolution of the density matrix, $\hat{\rho}(t) = e^{t\mathcal{L}}\hat{\rho}_0 = \langle\hat{U}_\eta(t)\hat{\rho}_0\hat{U}_\eta^\dagger(t)\rangle$, can then be recovered averaging over the fluctuations of the white noise, with $\hat{U}_\eta(t)$ the time evolution operator of the time-dependent Schrödinger equation at a fixed noise realization $\eta(t)$. The Hamiltonian Eq. (2) makes therefore clearer the connection of our setup to more conventional instances of OC, where algebraic scaling of the Loschmidt echo has been evidenced in quantum Ising models of the form Eq. (2) without adding a noisy character to the local perturbation [29,43].

However, since the state of the system is mixed at times $t > 0$, we need a generalized expression for the Loschmidt echo in order to investigate the onset of an analogue of OC in the dissipative critical quantum Ising chain. A natural choice is represented by the Uhlmann fidelity [44,45], which reduces to the Loschmidt echo when both states are pure. If instead only the initial state is pure (as in the case under inspection in this work), we observe that the Uhlmann fidelity retains a convenient expression

$$F(t) = \langle\psi(0)|\hat{\rho}(t)|\psi(0)\rangle = \mathrm{Tr}[\hat{\rho}(0)\hat{\rho}(t)], \tag{3}$$

which is amenable to analytical calculations. Intuitively, the Loschmidt echo for an open system is equivalent (within Born approximation) to the Uhlmann fidelity of a given subsystem if the environment remains unaffected during dynamics, since the latter can then be traced out [45] (provided the initial density matrix is a factorized product of the system and environment's density matrices).

At the critical point, the quantum Ising chain reacts to the presence of the local dephasing channel $\hat{L}$, with a fidelity which decays and scales at long times as

$$F(t) \propto t^{+\theta}e^{-\gamma t}, \qquad t \gg 1/J. \tag{4}$$

The power-law character $\propto t^{+\theta}$ recalls the characteristic algebraic response of a gapless quantum system to a local perturbation [2,3,29], which signals the onset of the phenomenon of orthogonality catastrophe. The exponent $\theta = 8/\pi^2(1-2n)^2(\kappa/J)^2$ is, however, *positive*, contrary to unitary incarnations of OC (n is the local fermion density on the site where the dissipative perturbation is applied, and it is a function of the transverse field g; see, for instance, Ref. [40]). This brings the qualitative difference that a new, concave region (see also Fig. 1 and the discussion in the following section) appears in the universal shape of $F(t)$, as a result of the interplay between $t^{+\theta}$ and the exponential decay $\propto e^{-\gamma t}$ with decoherence rate $\gamma = 8\kappa n(1-n)$—in contrast to the monotonic convex behavior of the Loschmidt echo in isolated systems. As in ordinary

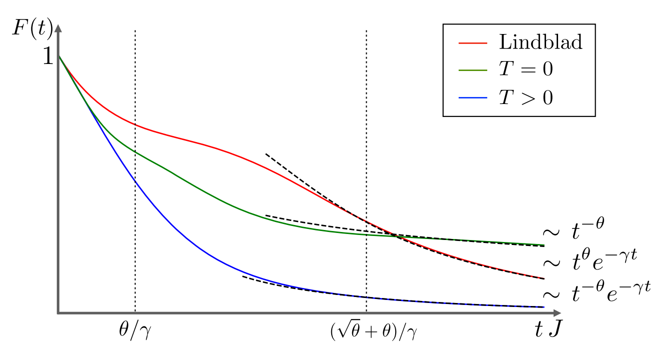

FIG. 1. Comparison between the fidelity $F(t)$ for a quantum Ising chain with a spin-dephasing impurity (red line) and for the same Ising chain with a local defect on the transverse field, at zero temperature [29] (green line) and at finite temperature [30] (blue line). The asymptotic behavior of the three curves is highlighted on the right. The local Lindbladian channel results in a slower decay of the fidelity compared to the other two cases. Close to $t \lesssim (\sqrt{\theta} + \theta)/\gamma$ the fidelity in the dissipative Ising model transits from a concave universal behavior to the usual convex character typical of isolated systems.

instances of orthogonality catastrophe, the power-law term is superseded when the many-body environment is away from criticality.

Cumulant expansion for Lindblad dynamics.—In order to find the long-time behavior Eq. (4), we design a second-order cumulant expansion for the fidelity suited for Lindblad dynamics, which generalizes analogous methods developed for the calculation of the Loschmidt echo in isolated systems [2,29]. The key idea is to express $F(t)$ in the superoperator formalism [46],

$$F(t) = \mathrm{Tr}[\hat{\rho}(0)e^{t\mathcal{L}}\hat{\rho}(0)] \equiv (\rho_0|e^{t\mathcal{L}}|\rho_0), \qquad (5)$$

where $e^{t\mathcal{L}}$ is the superoperator corresponding to the Lindblad dynamics in Eq. (1), acting on the supervector $|\rho_0)$ associated to the initial condition (the ground state of the quantum Ising chain in this specific instance). Casting $F(t)$ into the form Eq. (5) makes it amenable to a standard perturbative expansion in the interaction picture with respect to the unperturbed (purely Hamiltonian) Liouvillian $\mathcal{H}_0$, associated to the quantum dynamics of the Ising model. Within this representation, we evolve the density matrix, $\hat{\rho}_I(t) = e^{i\hat{H}_0 t}\hat{\rho}(0)e^{-i\hat{H}_0 t}$, starting from the critical ground state of the Ising chain, and we recast the fidelity using $\hat{\rho}_I(t)$ as the reference state,

$$F(t) = \left(\rho_I|T_{\leftarrow}\exp\left(+\int_0^t ds\mathcal{L}_I(s)\right)|\rho_I\right), \qquad (6)$$

which can then be expanded in cumulants (see Supplemental Material [47]),

$$F(t) = \exp\left(+\int_0^t ds[\mathcal{L}_I(s)]_0^C + \frac{1}{2}\int_0^t ds\int_0^t ds'[T_{\leftarrow}\mathcal{L}_I(s)\mathcal{L}_I(s')]_0^C + \cdots\right). \qquad (7)$$

In Eq. (7), $T_{\leftarrow}$ is the time ordering operator, $\mathcal{L}_I(s)$ is the Liouvillian perturbation with its Lindblad operators evolving under the Hamiltonian $\hat{H}_0$, we used $\hat{\rho}_I(t) = \hat{\rho}(0)$ for the initial ground state, and the compact notation $(\cdot)_0 \equiv (\rho_0|\cdot|\rho_0)$ has been adopted. For a single, Hermitian dissipative channel the first two cumulants read

$$[\mathcal{L}_I(s)]_0^C = -2(\langle\hat{L}^2\rangle_0 - \langle\hat{L}\rangle_0^2),$$
$$[T_{\leftarrow}\mathcal{L}_I(s)\mathcal{L}_I(s')]_0^C = 4[|\langle T_{\leftarrow}\hat{L}(s)\hat{L}(s')\rangle_0|^2 - \langle\hat{L}\rangle_0^4]. \qquad (8)$$

In order to gain insight into the first two terms of the cumulant expansion Eq. (8), we write them in terms of connected correlation functions of spin operators,

$$[\mathcal{L}_I(s)]_0^C = -2\kappa(1 - \langle\hat{\sigma}_j^z\rangle_0^2), \qquad (9)$$

$$[T_{\leftarrow}\mathcal{L}_I(s)\mathcal{L}_I(s')]_0^C = 8\langle\hat{\sigma}_j^z\rangle_0^2\mathrm{Re}G(s-s') + 4|G(s-s')|^2, \qquad (10)$$

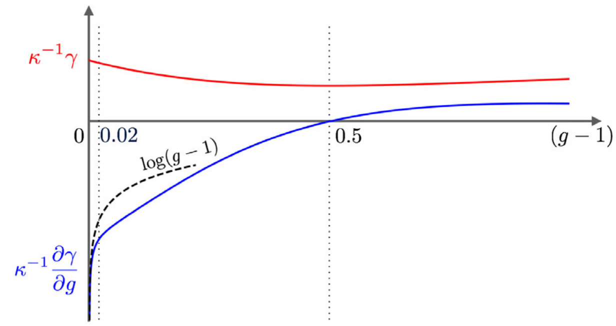

FIG. 2. The rate of exponential decay γ and its derivative as a function of the gap $(g-1)$ in the paramagnetic phase of the one-dimensional quantum Ising chain subject to local spin dephasing. Close to the critical point, $g \to 1$, the latter (blue line) exhibits a logarithmic divergence, while the former (red line) is continuous.

where $G(s) = \langle\hat{\sigma}_j^z(s)\hat{\sigma}_j^z(0)\rangle - \langle\hat{\sigma}_j^z\rangle_0^2$. The first cumulant Eq. (9) is constant, and when integrated over time yields a term proportional to t: this is the exponential decay rate γ in Eq. (4). γ is continuous close to the critical point $g \to 1$, where it has, however, a diverging derivative (see also Fig. 2):

$$\left.\frac{\partial\gamma}{\partial g}\right|_{g\to 1} = \frac{8-2\pi}{\pi^2}\kappa\log(g-1). \qquad (11)$$

This is a first imprint of criticality on the fidelity, although similar features have also been found in the study of decoherence induced on a two-level system coupled to a one-dimensional quantum spin chain [48].

The second cumulant Eq. (10) contains, instead, the characteristic features of the OC phenomenon; specifically, the first contribution to Eq. (10) diverges logarithmically in t after integration over the variables s and s' [cf. Eq. (7)]. Collecting Eq. (9) and this leading contribution, we have the following expression for the fidelity (see Supplemental Material for details [47]):

$$F(t) = \exp\left(-\gamma t + \kappa^2\int_0^t ds\int_0^t ds'\mathrm{Re}G(s-s') + \cdots\right)$$
$$= \exp\left(-\gamma t + 4\kappa^2(1-2n)^2\int_{k,k'}V(k,k') \times \frac{1-\cos(E_k+E_{k'})t}{(E_k+E_{k'})^2} + \cdots\right). \qquad (12)$$

In Eq. (12), $V(k,k') = \sin(2\theta_k)\sin(2\theta_{k'}) + 4\cos^2(\theta_k)\cos^2(\theta_{k'})$ is the same matrix element found in the second-order cumulant expansion of Ref. [29], with $2\theta_k = tg^{-1}[\sin k/(g-\cos k)]$. This logarithmic divergence is at the origin of the power-law character of Eq. (4), and it can be understood by power counting [the denominator is $\propto(k+k')^2$, V is finite, and integration over momenta k,k' is carried out twice]. A double time integration over a term

$\propto G(s-s')$ appears also in the second cumulant calculation of the Loschmidt echo in an isolated system, causing as well a logarithmic divergence in time and accordingly the typical algebraic scaling $\sim t^{-\theta'}$ [2,3,29,30]. The circumstance that the same quantity appears in the dissipative setup considered in this work at the same level of cumulant expansion confirms the physical intuition that also here critical correlations are the genuine cause of algebraic scaling. The same scaling argument shows that the term $\propto |G(s-s')|^2$ from Eq. (10) is subleading with respect to the terms appearing in Eq. (12).

We finally comment on the impact of the algebraic scaling $\propto t^{+\theta}$ in $F(t)$ [cf. Eq. (4)]. The fidelity is always monotonically decreasing, as it should be for a system coupled to a Markovian bath where there cannot be any revival of the information originally present in the initial state. $F(t)$ is apparently increasing for times $tJ \lesssim \theta J/\gamma \propto \kappa/J$, with $\kappa/J \ll 1$, the small parameter controlling the perturbative cumulant expansion; however, the algebraic scaling is only valid starting at times of the order $t \sim 1/J$ (as it occurs also in OC for isolated systems [29]), and therefore no actual growth occurs. Nevertheless, $F(t)$ displays a distinct feature compared to OC phenomena in closed systems: the presence of a gapless mode provokes the scaling $\propto t^{\theta}$ and decoherence is actually slowed down. Furthermore, the fidelity at early times is concave, see Fig. 1, and becomes convex at later times. The inflection point lies indeed at $t^* J = (\sqrt{\theta}+\theta)/\gamma$, which is $O(1)$ even for $\kappa/J \ll 1$. This behavior is general in the sense that it depends only on the long-time properties of the critical correlations of the model, and it constitutes a novelty of the dissipative scenario.

Other models.—We have tested the emergence of a dissipative analogue of OC in other systems, ranging from quantum spin chains with conserved local magnetization (XX model) to the two-dimensional Bose-Hubbard model with dephasing. We focused on the onset of the scaling term $\propto t^{+\theta}$ contributing to the fidelity, since aspects related to monotonicity and concavity are based on the generic structure of the perturbative cumulant expansion rather than on specific details of the model at hand.

The simplest generalization of the previous setup in one dimension is the XY spin chain [49] described by the Hamiltonian $\hat{H}_{XY} = -(J/2)\sum_i \{[(1+\Delta)/2]\hat{\sigma}_i^x\hat{\sigma}_{i+1}^x + [(1-\Delta)/2]\hat{\sigma}_i^y\hat{\sigma}_{i+1}^y + g\hat{\sigma}_i^z\}$, with a local dissipative impurity, $\hat{L} = \sigma_j^z$. For generic $\Delta \neq 0$, analytical results can be obtained from the cumulant expansion of the previous section simply replacing $\sin k \to \Delta \sin k$. The latter substitution does not alter the infrared scaling of Eq. (12), because the quasiparticle energy of the fermions diagonalizing H_{XY} still has a linear infrared character as $k \to 0$, $\epsilon_k \sim \Delta|k|$, implying that the fidelity has an algebraic scaling contribution also in this model. When $\Delta = 0$, the Hamiltonian H_{XY} describes an XX quantum spin chain [40], which conserves the total transverse magnetization ($\hat{M}_z \propto \sum_i \hat{\sigma}_i^z$); the model is therefore equivalent to a system of free fermions in one dimension at finite density, known to undergo orthogonality catastrophe when coupled to a local potential [2]. The dissipative analogue holds as well, the main difference with the Ising case being that the logarithmic divergence in Eq. (12) comes from modes close to the Fermi surface, rather than from those close to $k = 0$. In passing, this circumstance highlights that criticality is not a necessary condition for the onset of OC: the absence of a gap in the spectrum is sufficient to induce the long-range correlations that cause the algebraic scaling contribution to the fidelity.

Finally, we have considered the Bose-Hubbard model [40] in d spatial dimensions, $H_{\rm BH} = -J\sum_{\langle i,j\rangle}\hat{b}_i^\dagger\hat{b}_j + (U/2)\sum_i \hat{n}_i(\hat{n}_i - 1)$, deep in the superfluid phase (where excitations are gapless) and subject to a local heating process described by $\hat{L}_j = \hat{n}_j$, at rate κ; despite the fact that model is not at the critical point, the absence of a gap is sufficient to develop long-range correlations which make the model potentially prone to OC. In the Hartree-Fock-Bogolyubov approximation, the model reduces to a free Hamiltonian of Bogolyubov quasiparticles; computations follow the perturbative cumulant expansion Eq. (7) with the additional complication that now $\hat{L}^2 \neq \hat{\mathbf{1}}$, which brings a new term,

$$2\mathrm{Re}[\langle T_{\leftarrow}\hat{L}^2(s)\hat{L}^2(s')\rangle_0 - \langle\hat{L}^2\rangle_0^2], \qquad (13)$$

in Eq. (10). Employing scaling arguments, one can show that the phenomenon of dissipative OC exists only in $d = 2$, with a fidelity scaling as $F_{\rm BH}(t) \propto t^{\Theta}e^{-\Gamma t}$, where $\Gamma \propto \kappa n[1 + \mathcal{O}(n)]$, $\Theta \propto (\kappa/J)^2[1 + 4n + \mathcal{O}(n^2)]$ and n the density of bosons in the superfluid ground state. In passing, we notice that the interaction strength U determines the time scales, $t \gg (JUn)^{-1/2}$, for the onset of the scaling form $F_{\rm BH}(t)$ of the fidelity in the Bose-Hubbard model.

Conclusions and perspectives.—In summary, we have shown that the decoherence following the sudden switch of a dissipative impurity on a gapless quantum many-body system is slowed down due to the critical, long-range correlations persisting in the system. This phenomenon can be interpreted as another manifestation of the Anderson orthogonality catastrophe in the new context of driven-dissipative systems, thanks to the analogy to the stochastic quantum dynamics governed by the Hamiltonian Eq. (2). In fact, the potential is localized for every realization of the noise; hence, the transitions it can induce are suppressed in the low-frequency part of the spectrum as result of conventional orthogonality catastrophe physics. The corresponding absorption processes are inhibited in this energy window, and heating is therefore partially slowed, as explicated by the occurrence of a power-law growth t^{θ} together with the typical exponential decay $e^{-\gamma t}$.

A natural point to address is the transient nature of the phenomenon, i.e., whether dynamics is capable to exit the

OC regime at longer times due to heating. Therefore, as a future direction, we foresee a calculation of the fidelity for dissipative impurities with nonperturbative or numerical methods, in order to inspect whether its universal shape is a precursor of a pure relaxational regime entirely dominated by decoherence or whether it can persist for asymptotically long times (as it might happen in the context of quantum criticality in driven-dissipative platforms [50,51]).

A further option is represented by the extension of the present study to the case of a non-Markovian impurity or, equivalently, of a non-Markovian noisy transverse field [see Eq. (2) above]. In this scenario, a nonmonotonic behavior of the fidelity might be realizable due to the backflow of information from the environment to the system; accordingly, an intriguing possibility would be the existence of a time window where an algebraic growth is actually observable, unlike in the present case.

There is currently a research trend that aims at extending traditional topics in statistical mechanics to the domain of dissipative quantum many-body physics, as phase transitions [52] or integrability [53–55]. Our work articulates towards this direction; accordingly, a natural next step to substantiate the concept of a dissipative orthogonality catastrophe would consist in studying the response of driven-open fermionic or bosonic gapless systems [56–58] to local disturbances.

We thank E. Altman, A. Chiocchetta, H. Fröml, Z. Lenarcic, A. Rosch, and M. Schiró for interesting discussions on related topics. F. T. thanks L. Tolomeo for enlightening mathematical discussions. S. D. acknowledges funding by the German Research Foundation (DFG) through the Institutional Strategy of the University of Cologne within the German Excellence Initiative (ZUK 81) and SFB1238 project C04, and by the European Research Council via ERC Grant Agreement No. 647434 (DOQS). J. M. acknowledges support from the Alexander von Humboldt foundation. J. M. is supported by the European Union's Horizon 2020 research and innovation programme under the Marie Sklodowska-Curie Grant Agreement No. 745608 (QUAKE4PRELIMAT). This research was supported in part by the National Science Foundation under Grant No. NSF PHY-1748958.

[1] P. W. Anderson, Phys. Rev. Lett. **18**, 1049 (1967).
[2] G. Mahan, *Many-Particle Physics*, 3rd ed. (Springer, Berlin, 2000).
[3] P. Nozieres and C. T. De Dominicis, Phys. Rev. **178**, 1097 (1969).
[4] A. O. Gogolin, Phys. Rev. Lett. **71**, 2995 (1993).
[5] G. Yuval and P. W. Anderson, Phys. Rev. B **1**, 1522 (1970).
[6] P. W. Anderson, G. Yuval, and D. R. Hamann, Phys. Rev. B **1**, 4464 (1970).
[7] Y. Gefen, R. Berkovits, I. V. Lerner, and B. L. Altshuler, Phys. Rev. B **65**, 081106 (2002).
[8] C. Latta, F. Haupt, M. Hanl, A. Weichselbaum, M. Claassen, W. Wuester, P. Fallahi, S. Faelt, L. Glazman, J. von Delft, H. E. Türeci, and A. Imamoglu, Nature (London) **474**, 627 (2011).
[9] M. Knap, A. Shashi, Y. Nishida, A. Imambekov, D. A. Abanin, and E. Demler, Phys. Rev. X **2**, 041020 (2012).
[10] W. Münder, A. Weichselbaum, M. Goldstein, Y. Gefen, and J. von Delft, Phys. Rev. B **85**, 235104 (2012).
[11] B. Dóra, F. Pollmann, J. Fortágh, and G. Zaránd, Phys. Rev. Lett. **111**, 046402 (2013).
[12] A. Sindona, J. Goold, N. Lo Gullo, S. Lorenzo, and F. Plastina, Phys. Rev. Lett. **111**, 165303 (2013).
[13] R. Vasseur and J. E. Moore, Phys. Rev. B **92**, 054203 (2015).
[14] V. Khemani, R. Nandkishore, and S. L. Sondhi, Nat. Phys. **11**, 560 (2015).
[15] B. Dóra and F. Pollmann, Phys. Rev. Lett. **115**, 096403 (2015).
[16] R. Schmidt, M. Knap, D. A. Ivanov, J.-S. You, M. Cetina, and E. Demler, Rep. Prog. Phys. **81**, 024401 (2018).
[17] C. Weitenberg, M. Endres, J. F. Sherson, M. Cheneau, P. Schauss, T. Fukuhara, I. Bloch, and S. Kuhr, Nature (London) **471**, 319 (2011).
[18] T. Fukuhara, A. Kantian, M. Endres, M. Cheneau, P. Schauss, S. Hild, D. Bellem, U. Schollwck, T. Giamarchi, C. Gross, I. Bloch, and S. Kuhr, Nat. Phys. **9**, 235 (2013).
[19] A. Peres, Phys. Rev. A **30**, 1610 (1984).
[20] R. A. Jalabert and H. M. Pastawski, Phys. Rev. Lett. **86**, 2490 (2001).
[21] H. T. Quan, Z. Song, X. F. Liu, P. Zanardi, and C. P. Sun, Phys. Rev. Lett. **96**, 140604 (2006).
[22] A. Gambassi and A. Silva, arXiv:1106.2671.
[23] M. Heyl and S. Kehrein, Phys. Rev. B **85**, 155413 (2012).
[24] B. Dóra, A. Bácsi, and G. Zaránd, Phys. Rev. B **86**, 161109 (2012).
[25] R. Dorner, S. R. Clark, L. Heaney, R. Fazio, J. Goold, and V. Vedral, Phys. Rev. Lett. **110**, 230601 (2013).
[26] L. Mazzola, G. De Chiara, and M. Paternostro, Phys. Rev. Lett. **110**, 230602 (2013).
[27] C. Karrasch, S. Andergassen, M. Pletyukhov, D. Schuricht, L. Borda, V. Meden, and H. Schoeller, Europhys. Lett. **90**, 30003 (2010).
[28] M. Heyl, A. Polkovnikov, and S. Kehrein, Phys. Rev. Lett. **110**, 135704 (2013).
[29] A. Silva, Phys. Rev. Lett. **101**, 120603 (2008).
[30] C. Lupo and M. Schiró, Phys. Rev. B **94**, 014310 (2016).
[31] M. Schiró and A. Mitra, Phys. Rev. Lett. **112**, 246401 (2014).
[32] T. Gericke, P. Würtz, D. Reitz, T. Langen, and H. Ott, Nat. Phys. **4**, 949 (2008).
[33] T. Gericke, J. Meineke, D. Stadler, S. Krinner, and T. Esslinger, Science **337**, 1069 (2012).
[34] V. A. Brazhnyi, V. V. Konotop, V. M. Pérez-García, and H. Ott, Phys. Rev. Lett. **102**, 144101 (2009).
[35] P. Würtz, T. Langen, T. Gericke, A. Koglbauer, and H. Ott, Phys. Rev. Lett. **103**, 080404 (2009).
[36] D. A. Zezyulin, V. V. Konotop, G. Barontini, and H. Ott, Phys. Rev. Lett. **109**, 020405 (2012).
[37] G. Barontini, R. Labouvie, F. Stubenrauch, A. Vogler, V. Guarrera, and H. Ott, Phys. Rev. Lett. **110**, 035302 (2013).

[38] T. T. S. Nakajima, I. Danshita, Y. Takasu, and Y. Takahashi, Sci Adv. **3**, e1701513 (2017).

[39] H. Froml, A. Chiocchetta, C. Kollath, and S. Diehl, arXiv:1809.09085 [Phys. Rev. Lett. (to be published)].

[40] S. Sachdev, *Quantum Phase Transitions* (Cambridge University Press, Cambridge, England, 2011).

[41] J. Marino and A. Silva, Phys. Rev. B **89**, 024303 (2014).

[42] K. Stannigel, P. Hauke, D. Marcos, M. Hafezi, S. Diehl, M. Dalmonte, and P. Zoller, Phys. Rev. Lett. **112**, 120406 (2014).

[43] P. Smacchia and A. Silva, Phys. Rev. Lett. **109**, 037202 (2012).

[44] A. Uhlmann, Rep. Math. Phys. **9**, 273 (1976).

[45] M. A. Nielsen and I. L. Chuang, *Quantum Computation and Quantum Information* (Cambridge University Press, Cambridge, England, 2012).

[46] H.-P. Breuer and F. Petruccione, *The Theory of Open Quantum Systems* (Oxford University Press, Oxford, 2007).

[47] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.122.040604 for details on the cumulant expansion for the fidelity.

[48] D. Rossini, T. Calarco, V. Giovannetti, S. Montangero, and R. Fazio, Phys. Rev. A **75**, 032333 (2007).

[49] E. Lieb, T. Schultz, and D. Mattis, Ann. Phys. (N.Y.) **16**, 407 (1961).

[50] E. G. Dalla Torre, E. Demler, T. Giamarchi, and E. Altman, Phys. Rev. B **85**, 184302 (2012).

[51] J. Marino and S. Diehl, Phys. Rev. B **94**, 085150 (2016).

[52] L. M. Sieberer, M. Buchhold, and S. Diehl, Rep. Prog. Phys. **79**, 096001 (2016).

[53] L. Banchi, D. Burgarth, and M. J. Kastoryano, Phys. Rev. X **7**, 041015 (2017).

[54] M. V. Medvedyeva, F. H. L. Essler, and T. Prosen, Phys. Rev. Lett. **117**, 137202 (2016).

[55] D. A. Rowlands and A. Lamacraft, Phys. Rev. Lett. **120**, 090401 (2018).

[56] J. Eisert and T. Prosen, arXiv:1012.5013.

[57] M. Höning, M. Moos, and M. Fleischhauer, Phys. Rev. A **86**, 013606 (2012).

[58] B. Horstmann, J. I. Cirac, and G. Giedke, Phys. Rev. A **87**, 012108 (2013).