

Topological and nontopological defects from quantum reset dynamics

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We analyze mechanisms for universal out-of-equilibrium dynamics near criticality by exploring the effect of randomized quantum resetting (QR) under a linear finite-time drive across, or very close to, an equilibrium quantum critical point. Using the transverse-field Ising chain as a minimal model and exploiting its exact solution, QR is found to cause a crossover of the scaling of the defect density with the time scale τ of the drive, from Kibble-Zurek to anti-Kibble-Zurek scaling as τ increases. Depending on whether the ramp ends in a symmetry-broken phase or not, the resulting defects are topological or nontopological. The competition between QR and the Kibble-Zurek mechanism gives rise to local minima of the defect densities at optimal annealing times. These times and the corresponding local minima are shown to scale as universal power laws with the rate of QR. Additional results for the scaling of the mean excess energy suggest that a system driven across, or very close to, a quantum critical point exhibits the same scaling behavior under a linear finite-time drive with QR as with uncorrelated noise. We suggest that these results may serve as benchmarks for validating quantum annealing devices.

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I. INTRODUCTION

The study of nonequilibrium quantum matter has surged in the last two decades, driven by experimental advances across various platforms such as ultracold gases [1], Rydberg atom arrays [2], digital NISQ [3] and analog [4] quantum computing devices, in addition to a growing class of different quantum materials [5]. These developments have raised a number of theoretical challenges, including the quest for universal mechanisms that govern the out-of-equilibrium dynamics of quantum many-body systems [6].

One of the few established paradigms for universality away from equilibrium is the Kibble-Zurek (KZ) mechanism [7,8], extended in Refs. [9–13] to the quantum realm: A system driven into a broken-symmetry phase by varying a control parameter across the critical point of a continuous equilibrium phase transition is predicted to violate adiabaticity due to critical slowing down, leading to the appearance of topological defects. At the core of the KZ mechanism is a universal scaling law $n \sim \tau^{-\beta}$ that governs how the average density of defects n depends on the time scale $\tau = |dt/d\lambda|$ (inverse driving rate) of a linear ramp $\lambda(t) = \text{const.} \pm t/\tau$ across the

critical point, with exponent $\beta = d\nu/(1+z\nu)$ where d is the dimensionality, and z and ν the equilibrium dynamic and correlation length critical exponents, respectively [8].

A large number of experiments on near-isolated quantum systems show support for KZ scaling (see Ref. [14] and references therein). However, environmental disturbances can significantly influence the dynamics and cause its breakdown as first reported in an experiment on a multiferroic [15]. This experiment found an increase of the average defect density n at large time scales, later found to be characteristic of anti-KZ behavior where n scales linearly with τ , $n \sim \kappa\tau$, with κ the rate by which noise from the environment induces excitations in the system [16].

Noise-induced anti-KZ scaling is by now well studied [16–25]. As for KZ scaling, it serves as a useful testbed for validating implementations of quantum annealing schemes [26–31]—key to simulation and optimization protocols for analog quantum computing [32,33]. The nonadiabatic excitations that come with a finite-time ramp which crosses a critical point and underlies the KZ mechanism are to be addressed by “shortcuts to adiabaticity” [34] or other control strategies [35], necessary for flawless adiabatic quantum computing when used for optimization tasks.

Noise is not the only way to induce stochastic behavior in a system, so it would be interesting to understand if and how anti-KZ emerges generally with nonequilibrium randomness, e.g. by using resetting protocols (for a review see Ref. [36]). In this article we propose that adding so-called *quantum resetting* (QR) [37–45] to a finite-time ramp across, or very close to, a quantum critical point may serve as an additional useful and precise test of a quantum annealing device. As we shall

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show, when a time-evolved state is reset to its initial state at random times during the ramp, this systematically induces a crossover from KZ scaling with exponent β to linear anti-KZ scaling, but with no noise present. The defect production is that of topological or nontopological defects, depending on whether the ramp ends in a symmetry-broken phase or not. Moreover, the behavior shows data collapse and is described by universal critical exponents, indicating that the crossover is robust and predictable and can be used as a baseline for evaluating the performance of a quantum annealing process. In contrast, when the ramp is located within a phase away from the critical point, the defect production gets drastically reduced and only some features of scaling survive, with others being corrupted.

QR is typically realized by stochastic resets of the driving of a closed system [37–45]. The time-evolved density matrix in this way gets reset to its initial condition at random times, characterized by some reset rate r . QR is a quantum descendant of reset dynamics known to accelerate search processes and induce nonequilibrium steady states in classical systems (see Ref. [36] and references therein). The latter property is inherited by QR dynamics where the unitary time evolution of a closed quantum system repeatedly interrupted by stochastic resets asymptotically results in a steady-state mixed density matrix, exhibiting nonzero off-diagonal elements [37]. Additional features uncovered by studies of quantum walks with QR [38,40,43], open system dynamics resulting from QR [39,42], and entanglement generation from QR [44,45], suggest QR as a tool to achieve nonequilibrium steady states with certain specified properties [41].

Here, we shall focus on the effect of QR on the generation of topological and nontopological defects in a transverse-field Ising (TFI) chain when driven by a linear ramp across one of its equilibrium quantum critical points. The TFI chain serves as a minimal and paradigmatic model for the study of quantum phase transitions [46] and KZ physics [13,14]. The simplicity of the model, its rich nonequilibrium spin dynamics, and the availability of an exact analytical solution [47], has made it useful for benchmarking simulations on various D-Wave computers [26–30]—with such architectures expected to allow for QR [48–50].

II. PRELIMINARIES

A. Transverse-field Ising chain

The Hamiltonian of the TFI chain with a time-dependent magnetic field can be written as

$$H(t) = - \sum_{n=1}^N (\sigma_n^x \sigma_{n+1}^x + h(t) \sigma_n^z), \quad (1)$$

setting the overall energy scale to unity and choosing N even, with periodic boundary conditions on the Pauli matrices, $\bar{\sigma}_1 = \bar{\sigma}_{N+1}$. We shall take $h(t) = h_i \pm t/\tau$ to be a linear ramp, from an initial value h_i at $t=0$ to a final value h_f at $t=t_f$. Recall that when $h = \pm 1$, the model exhibits quantum phase transitions in the thermodynamic limit, with paramagnetic phases for $|h| > 1$ and a ferromagnetic phase for $|h| < 1$ [47].

The instantaneous eigenstates of the TFI Hamiltonian $H(t)$ serve as useful basis states for the analysis of the

nonequilibrium dynamics. These states are obtained via a Jordan-Wigner transformation to fermionic creation- and annihilation operators $c_j^\dagger$ and c_j [51] whereby the transformed TFI Hamiltonian becomes block diagonal, $H(t) = H_e(t) \otimes H_o(t)$, with $H_{e/o}$ acting on states in the Fock subspace with an even/odd number of fermions. Choosing to work in the even subspace where the fermions obey antiperiodic boundary conditions, $c_{N+1} = -c_1$, the Fourier transformation $c_j = N^{-1/2} \sum_k e^{i(jk-\pi/4)} c_k$ turns $H_e(t)$ into a sum over decoupled first-quantized mode Hamiltonians $\mathcal{H}_k(t)$,

$$H_e(t) = \sum_{k>0} C_k^\dagger \mathcal{H}_k(t) C_k. \quad (2)$$

Here, $C_k^\dagger = (c_k^\dagger \ c_{-k})$ are Nambu spinors and $\mathcal{H}_k(t) = h_k(t) \sigma^z + \Delta_k \sigma^x$ where $h_k(t) = 2(h(t) - \cos(k))$ and $\Delta_k = 2 \sin(k)$, with $k \in [\pm\pi/N, \pm 3\pi/N, \dots, \pm(N-1)\pi/N]$ as implied by the antiperiodic boundary conditions [52].

An instantaneous Bogoliubov transformation $\gamma_k(t) = \cos \theta_k(t) c_k - \sin \theta_k(t) c_{-k}^\dagger$ diagonalizes $H_e(t)$,

$$H_e(t) = \sum_k \varepsilon_k(t) \left(\gamma_k^\dagger(t) \gamma_k(t) - \frac{1}{2} \right), \quad (3)$$

with $\theta_k(t) = \arctan(\Delta_k/[h_k(t) - \varepsilon_k(t)])$ and $\varepsilon_k(t) = \sqrt{h_k^2(t) + \Delta_k^2}$. The eigenstates of $H_e(t)$ are

$$\begin{aligned} |\phi_k^-(t)\rangle &= (\cos \theta_k(t) + \sin \theta_k(t) c_k^\dagger c_{-k}^\dagger) |0\rangle \\ |\phi_k^+(t)\rangle &= (-\sin \theta_k(t) + \cos \theta_k(t) c_k^\dagger c_{-k}^\dagger) |0\rangle, \end{aligned} \quad (4)$$

with energies $\varepsilon_k^\pm(t) = \mp \varepsilon_k(t)$, and where $|0\rangle$ is the vacuum of $\{c_k, c_k^\dagger\}$. By repeating the procedure above for $H_o(t)$, one establishes that the instantaneous ground state of the full TFI Hamiltonian $H(t)$ lies in the even subspace for finite N , allowing us to initialize the ramp in the ground state of the system.

B. Topological defects from ramping without QR

A topological defect in the TFI chain takes the form of a domain wall (or “kink”) in the ferromagnetic phase $|h| < 1$, being a boundary between regions with different spin orientations [46]. These defects are delocalized and in 1-1 correspondence with the quasiparticle excitations $\gamma_k^\dagger |\phi_k^-(t)\rangle$ [52], appearing pairwise in the even subspace and having energies $\varepsilon_k(t)$ according to Eq. (3). It follows that the expected number of defects $\mathcal{N}_0$ at the end of a ramp, from the paramagnetic into the ferromagnetic phase, is given by $\mathcal{N}_0 = \sum_k^N \langle \psi_k(t_f) | \gamma_k^\dagger(t_f) \gamma_k(t_f) | \psi_k(t_f) \rangle$ for large N , with $|\psi_k(t)\rangle$ being the state of mode k at time t ,

$$|\psi_k(t)\rangle = \alpha_k^+(t) |\phi_k^+(t)\rangle + \alpha_k^-(t) |\phi_k^-(t)\rangle. \quad (5)$$

Note that while the amplitudes $\alpha_k^\pm(t)$ in Eq. (5) are time evolved from the initial state, $|\phi_k^\pm(t)\rangle$ are instantaneous.

It can be verified that $\gamma_k^\dagger(t) \gamma_k(t) |\phi_k^+(t)\rangle = |\phi_k^+(t)\rangle$ and $\gamma_k^\dagger(t) \gamma_k(t) |\phi_k^-(t)\rangle = 0$, so it follows from the above that the expected number of defects is given by

$$\mathcal{N}_0 = \sum_k \langle \phi_k^+(t_f) | \rho_{0,k}(t_f) | \phi_k^+(t_f) \rangle, \quad (6)$$

where $\rho_{0,k}(t) = |\psi_k(t)\rangle \langle \psi_k(t)|$ is the density matrix in the $\{|\phi_k^\pm(t)\rangle\}$ basis. By exploiting that the eigenstates of $\mathcal{H}_k(t)$ are

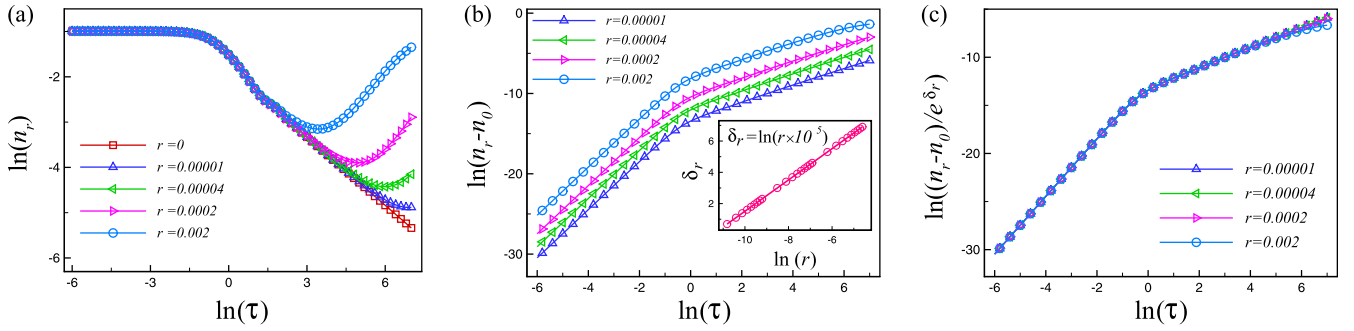


FIG. 1. (a) Log-log plot of the density of defects n_r generated by a ramp with QR from $h_i(0) = 2$ to $h_f(2\tau) = 0$ as a function of the time scale τ of the drive for different values of QR rate r . (b) Difference $\delta n_r = n_r - n_0$ between the density of generated defects from the ramp in (a) with $[n_r]$ and without $[n_0]$ QR. (c) Data collapse of the curves in panel (b) after rescaling $\delta n_r \rightarrow e^{-\delta_r} \delta n_r$. All data in the figure are obtained from a chain with 1000 sites.

the first-quantized equivalents of $|\phi_k^\pm(t)\rangle$, realized by the mapping $|0\rangle \rightarrow (01)^T$, $c_k^\dagger c_{-k}^\dagger |0\rangle \rightarrow (10)^T$ in Eq. (4), we solve the von Neumann equation for $\rho_{0,k}(t)$ exactly using a result by Vitanov and Garraway for Landau-Zener-type problems [53]. This yields the expected KZ scaling $\mathcal{N}_0 \sim \tau^{-1/2}$, first derived for the TFI chain by Dziarmaga via a different route [12].

III. TOPOLOGICAL DEFECTS FROM RAMPING WITH QR

A. Formalism

When ramping with QR, the unitary time evolution gets interrupted at random times by resets to the initial state at time $t=0$. We shall focus on ideal resets $|\psi(t)\rangle \rightarrow |\psi(0)\rangle$ where the projection to the initial state $|\psi(0)\rangle = \otimes_k |\psi_k(0)\rangle$ has unit probability. Moreover, as is standard, the distribution of random resets will be taken as Poissonian. As follows from the mode decomposition, $|\psi(t)\rangle = \otimes_k |\psi_k(t)\rangle$ together with Eq. (2), the QR protocol for $|\psi(t)\rangle \rightarrow |\psi(0)\rangle$ that enforces these two conditions can formally be broken down to

$$|\psi_k(t+dt)\rangle = \begin{cases} |\psi_k(0)\rangle, & \mathcal{P} = rdt \\ (1 - i\mathcal{H}_k(t)dt) |\psi_k(t)\rangle, & \mathcal{P} = 1 - rdt, \end{cases} \quad (7)$$

where $\mathcal{P}$ denotes the probability and r is the rate of resets. Hence, within a time interval dt , mode k is reset to its initial state with probability rdt , or, with the complementary probability $(1 - rdt)$, the mode is time-evolved by $\mathcal{H}_k(t)$. With $\mathcal{H}_k(t)$ being the first-quantized Hamiltonian, note that the states in the protocol are the first-quantized equivalents of the states in Eq. (5) obtained as above.

Guided by Eq. (6) for a ramp without QR, the problem to determine the expected number of defects *with* QR boils down to computing the corresponding density matrix. While the nonunitary time evolution for a fixed single realization of time instants for QR events, call it η , is deterministic, the stochastic QR protocol in Eq. (7) introduces an element of classical randomness in the sense that the time instants for QR events vary from one realization to another. The mixed density matrix $\rho_{r,k}(t)$ in the presence of QR is therefore obtained as the expectation value $\rho_{r,k}(t) = \langle \rho_{\eta,k}(t) \rangle_r$ by averaging over all

QR realizations $\{\eta\}$ with rate r . As a result [37],

$$\rho_{r,k}(t) = r \int_0^t e^{-rt'} \rho_{0,k}(t') dt' + e^{-rt} \rho_{0,k}(t), \quad (8)$$

where the last term corresponds to the event that there is no QR within $t \in [0, t]$.

By substituting $\rho_{r,k}(t_f)$ for $\rho_{0,k}(t_f)$ in Eq. (6), we can numerically obtain the expected number of defects $\mathcal{N}_r$ in a ramp with QR of rate r .

B. Ramps from the paramagnetic to the ferromagnetic phase

For definiteness we shall assume that the system is initially $[t=0]$ prepared in its paramagnetic ground state $|\psi(0)\rangle = \otimes_k |\phi_k^-(0)\rangle$ with $h_i = 2$. The magnetic field is then ramped through the quantum critical point $h = 1$ to the final value $h_f = 0$ in the ferromagnetic region at $t_f = 2\tau$.

The dependence of the expected density of defects $n_r = \mathcal{N}_r/N$ on the time scale τ of the ramp is shown in Fig. 1(a) for several values of the QR rate r . For $r=0$ (no QR), we numerically verify the known KZ scaling law for the TFI chain, $n_0 \sim \tau^{-1/2}$ [12] when $\tau \gtrsim 1$. The plateau that develops for smaller values of τ reflects the saturation of defects for very fast ramps where KZ scaling breaks down [54]. Given the scale in Fig. 1(a), the effect of QR ($r \neq 0$) is invisible on the plateau as well as in the adjacent r -dependent τ -intervals where KZ scaling still holds. Beyond these intervals, for larger values of τ , the defect densities n_r start to deviate from KZ scaling. This is indicative of a crossover to an anti-KZ regime [16,17] where slower ramps boost the defect density. As τ increases further, n_r passes through a local minimum and then shoots up, characteristic of anti-KZ. Thus, the defect generation appears to be controlled by two competing processes: (i) ramp-driven quasiparticle excitations, suppressed by increasing τ as in a KZ scenario [12], and (ii) QR-induced quasiparticle excitations, amplified by increasing τ as in an anti-KZ scenario [16,17]. Likewise, smaller values of τ lead to higher nonadiabaticity and provide less opportunity for QR to prevail. As manifest in Fig. 1(a), the competition between the two processes yield local minima of n_r at time scales $\tau_{\text{opt},r}$, acting as optimal times in a putative quantum annealing scheme [32], to be discussed below.

The QR-induced excitations can heuristically be understood as coming from the effective sudden quench experienced by the state $|\psi(0)\rangle$ right after resetting $|\psi(t_r)\rangle \rightarrow |\psi(0)\rangle$ at some time t_r . Different from the initial condition, the state $|\psi(0)\rangle$ does not see a ramp that smoothly takes off from $h(0)=0$, but instead a ramp that starts from a finite value $h(t_r)$, thus appearing as a sudden quench at t_r . Since sudden quenches boost the number of excitations [55], a slow ramp (large τ) that allows for many QRs will enable these to produce a larger number of excitations than a fast ramp (in contrast to the KZ mechanism that yields fewer excitations as τ increases).

Figure 1(b) displays scaling of the defect density from QR only, obtained by subtracting n_0 from n_r . Extracting numbers from the data, QR-induced defects are found to exhibit anti-Kibble-Zurek scaling with a linear growth with τ , $\delta n_r \sim \tau^\alpha$ where $\alpha = 1.00 \pm 0.02$ when $\tau > 1$. In contrast, when $\tau < 1$, the growth of defects with τ scales nonlinearly with an exponent $\alpha' = 3.000 \pm 0.003$. The sharp jump of the exponent at $\tau = 1$ —which happens before the density-saturated plateau has fully developed—is suggestive of some type of criticality in the QR dynamics. As transpires from Fig. 1(b), the number of defects when $\tau < 1$ is likely too small to be seen in an experiment, but from a theoretical point of view the manifest scaling with a well-defined exponent α' points to an interesting underlying mechanism. Let us add that finite-size effects for the scaling displayed in Fig. 1(b) for $N = 1000$ are negligible, while becoming visible—in the range of a few percent—when $N \approx 100$ but then only for very slow ramps with $\ln(\tau) \gtrsim 6$.

As seen in Fig. 1(b), the δn_r curves undergo displacements as the value of r varies. An analysis reveals that the shift for a given value of r , call it δ_r , scales logarithmically with r as $\delta_r = \ln(r \times 10^5)$ [inset of Fig. 1(b)]. This finding implies that the scaling of $e^{-\delta_r} \delta n_r$ with τ for different values of r is universal. Figure 1(c) illustrates the universality, showing that all curves from Fig. 1(b) collapse onto a single graph after the rescaling.

Summing up the data in Fig. 1(a), with input from Fig. 1(b), the defect density n_r behaves as

$$n_r \approx \begin{cases} h(r)\tau^{\alpha'} + \Lambda(\tau), & \tau < 1 \\ h(r)\tau^\alpha + b\tau^{-\beta}, & \tau > 1 \end{cases}, \quad (9)$$

with $\alpha' = 3.000 \pm 0.003$, $\alpha = 1.00 \pm 0.02$, $\beta = 0.50 \pm 0.01$. The amplitude $h(r)$ is the r -dependent rate at which QR generates quasiparticle excitations, with $\Lambda(\tau)$ a function that describes the crossover to the density-saturated plateau [56] in Fig. 1(a) as τ decreases below unity. Here, $b = c^{-1/2}$, where c is the speed of quasiparticle excitations [10], of order unity for $\hbar = a = 1$, with a the lattice spacing of the TFI chain. Note that the values of the exponents α and α' are extracted in the rate interval $r \in [0, 2 \times 10^{-2}]$, as indicated in Fig. 1. Pushing the QR rate beyond this interval will introduce rapidly increasing error bars for the critical exponents, eventually resulting in a breakdown of scaling when $r \gtrsim 0.02$. This is analogous to noise-induced anti-KZ scaling which assumes weak-intensity noise [16].

Apart from the QR-dependent prefactor $h(r)$ and the specific error estimates of the exponents in Eq. (9), the behavior

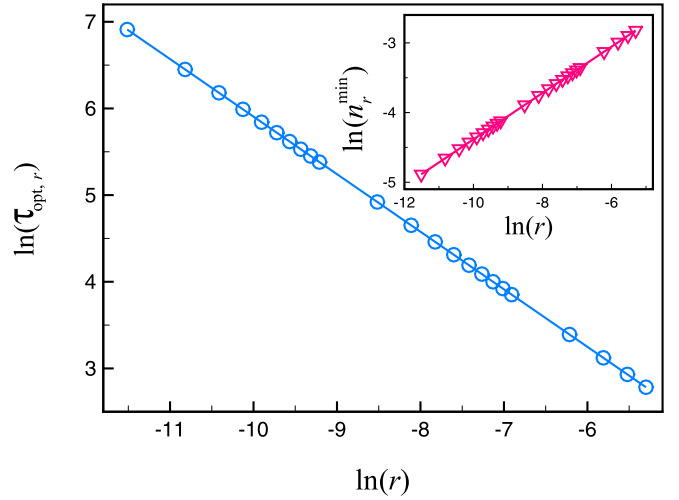


FIG. 2. Scaling of the optimal quantum annealing times $\tau_{\text{opt},r}$ with QR rate r for a ramp from $h_i(0) = 2$ to $h_f(2\tau) = 0$. Inset: Scaling of the local minima of defect densities n_r^{min} with r .

of n_r for $\tau > 1$ is the same as that obtained by adding uncorrelated (white) noise [16] or correlated (colored) fast noise [17,18] to a ramp across criticality of the TFI chain. The identity of scaling exponents for the two types of ramp dynamics—extracted in the limits of low-rate QR and weak-intensity noise, respectively—is most intriguing, specifically since the probability distributions in their ramp protocols, Poissonian for QR and Gaussian for noise, are different. This invites the question if this similarity also extends to other observables, such as the time $\tau_{\text{opt},r}$ of minimal defects introduced above.

Optimal times for quantum annealing define the best trade-offs between maximal adiabaticity (slow ramp) and minimal decoherence (fast ramp), at the same time constraining the defect production when crossing a critical point [32]. In the presence of QR, the location of the local minima of the r -dependent graphs in Fig. 1 pinpoint distinct optimal times $\tau_{\text{opt},r}$ beyond which there is a fast increase of defects. As follows from the data shown in Fig. 2, $\tau_{\text{opt},r} \sim r^{-\gamma}$ with $\gamma = 0.664 \pm 0.002 \simeq 2/3$. Also this scaling exhibits a perfect match between QR and noisy ramp dynamics: The exponent γ is the same as that which governs the scaling of the optimal annealing time with noise intensity in a standard quantum annealing protocol with white [16] or fast colored [17,18] weak-intensity noise. Our numerical results further show that the local minima of the defect densities n_r^{min} scale with the QR rate as $n_r^{\text{min}} \sim r^{0.332 \pm 0.002} \simeq 1/3 = \gamma/2$ (inset of Fig. 2). The fact that the scaling of defect densities and optimal annealing times in QR dynamics are identical with those in ramp dynamics with fast noise (uncorrelated or correlated) makes it compelling to analyze additional quantities. One such object, computed also in Ref. [16], is the mean excess energy of mode k with respect to the corresponding ground state at the end of the ramp [57],

$$Q_r = \frac{1}{N} (\text{Tr}(\rho_{r,k}(t_f) \mathcal{H}_k(t_f)) - \langle \phi_k^-(t_f) | \mathcal{H}_k(t_f) | \phi_k^-(t_f) \rangle). \quad (10)$$

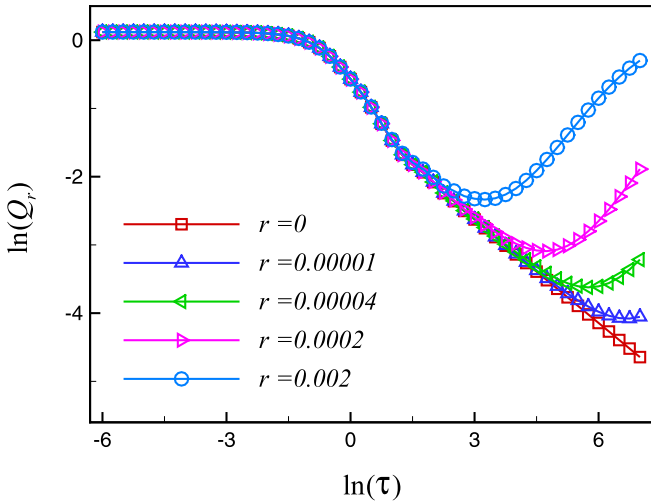


FIG. 3. Mean excess energy Q_r after a ramp from $h_i(0) = 2$ to $h_f(2\tau) = 0$ with QR for different rates r as a function of the time scale τ of the drive.

As suggested by a comparison between Figs. 1(a) and 3, one finds that Q_r exhibits the same universal scaling as the defect density in Eq. (9). This result agrees with the noisy ramp dynamics in Ref. [16] which shows identical scaling behaviors for defect densities and excess energies. The universality between the dynamics enforced by the two ramp protocols, one with QR and the other with noise, extends to the scaling of the variance of the mean excess energies, as well as to the optimal annealing times extracted from the local minima of the excess energies for different rates r (cf. Fig. 3).

IV. NONTOPOLOGICAL DEFECTS FROM RAMPING WITH QR

A. Ramps from the ferromagnetic to the paramagnetic phase

Figure 1(a) suggests that the KZ mechanism dominates the scaling of the topological defect density in the interval $1 < \tau < \tau_{\text{opt},r}$, within which increasing time scales imply a decreasing density. Recall that the KZ mechanism rests on the assumption that the critical slowing down in the vicinity of a critical point enforces local selections of a spontaneously broken symmetry, in this way producing topological defects that separate regions with different order [14]. However, we work in the even-parity sector of the fermionized TFI chain with a unique ground state, hence with no spontaneous symmetry breaking in the ferromagnetic region. Still, the KZ exponent $\beta = 1/2$ known for the TFI chain is well reproduced by our data. This is so, since a quasiparticle excitation in the fermionized TFI chain signals the presence of a region of mismatch between spins—costly in energy—independent of the underlying type of order or its origin.

This becomes manifest when ramping from the ferromagnetic to the paramagnetic phase of the TFI chain where the defects are nontopological, realized by single spin flips [46]. Here, KZ scaling applies also after a ramp from $|h_i| < 1$ (ferromagnet) to $h_f > 1$ (paramagnet) with no spontaneous symmetry breaking in the paramagnetic phase [12]. It follows that KZ scaling has a wider applicability than predicted by

the standard KZ *mechanism* (based on spontaneous symmetry breaking and an impulse approximation scenario [14]). Considering QR added to a ramp from a ferromagnet to a paramagnet, using the same mode-decomposed QR protocol as in Eq. (7) but with $h_i = 0$ at $t = 0$ and $h_f = 2$ at $t_f = 2\tau$, we find that all scaling exponents α , α' , β , and γ reported in Sec. III B are well reproduced, with roughly the same numerical errors. The only change is that the optimal annealing times $\tau_{\text{opt},r}$, signaling the onset of the anti-KZ regimes, are shifted by a small amount compared to Fig. 1(a).

B. Ramps within the paramagnetic phase

Recent results by one of us [58]—showing that KZ scaling in some models may persist for a ramp outside a quantum critical point—raises the question: What happens if adding QR to a noncrossing ramp in the TFI chain?

Focusing on the paramagnetic phase, Fig. 4 displays three cases: the defect density generated by (a) a ramp that terminates right before the critical point $h_c = 1$; (b) a ramp that terminates at an intermediate distance from h_c ; and (c) a ramp that terminates far away from h_c . These three panels illustrate the generic behavior for the TFI universality class.

When $r = 0$ (no QR), and using the same scaling analysis as in Sec. III B, we find that the KZ exponent $\beta = 1/2$ is modified very little (only in the second decimal point) when the ramp terminates right before the critical point. As one moves away from the critical point there is a rapid shift to a new exponent $\beta' = 2 \pm 0.03$, governing the $r=0$ scaling (cf. panels (b) and (c) of Fig. 4). The region of this KZ-like scaling, $n_0 \sim \tau^{-\beta'}$, shrinks with the distance from criticality as does the expected defect density n_0 for $\tau \approx 0$ (sudden quench). The latter feature reflects the fact that the deeper one goes with the ramp into the paramagnetic phase, the more difficult it is to overcome the gap to the first instantaneous excited state (with the gap increasing with the size of the magnetic field). While likely visible in a quantum annealing simulation, the extremely small number of expected nontopological defects away from criticality makes it improbable that they would be detected in a real experiment. This may explain why the $r=0$ scaling with exponent $\beta' \approx 2$ has gone unnoticed or has been considered irrelevant.

Turning to the QR-induced crossover ($r \neq 0$) to a regime with an increase of the expected defect density n_r with τ , one verifies—again by the same analysis as in Sec. III B—that the scaling of $n_r - n_0$ in all three cases in Fig. 4 here remains linear, i.e., with an exponent $\alpha = 1$, characteristic of anti-KZ scaling. When the ramp terminates close to, or at an intermediate distance from, criticality (cf. panels (a) and (b) in Fig. 4), this scaling sets in at values of τ before the crossover, while further away from criticality [cf. panel (c)] the linear scaling of $n_r - n_0$ appears approximately simultaneous with the crossover. Only for a ramp that terminates very close to the critical point does the optimal annealing time $\tau_{\text{opt},r}$ and mean excess energy Q_r exhibit universal scaling, with an exponent that differs slightly (in the second decimal) from that for a ramp crossing quantum criticality (cf. Secs. III B and IV B).

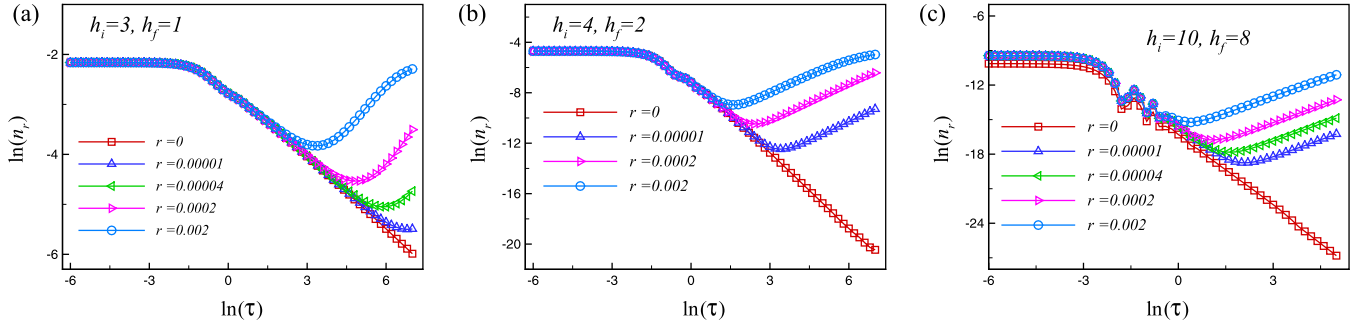


FIG. 4. (a) Log-log plot of the density of defects n_r generated by a ramp with QR from $h_i(0) = 3$ to $h_f(2\tau) = 1$ as a function of the time scale τ of the drive for different values of QR rate r . (b) Ditto, from $h_i(0) = 4$ to $h_f(2\tau) = 2$. (c) Ditto, from $h_i(0) = 10$ to $h_f(2\tau) = 8$. All data in the figure are obtained from a chain with 1000 sites.

V. SUMMARY

The universal behavior of linear ramps across, or very close to, an equilibrium quantum critical point has been analyzed in the presence of QR, using the TFI chain as a minimal model. The out-of-equilibrium dynamics near criticality with QR clearly shows the emergence of anti-KZ behavior, notably with data collapse and scaling. In particular, the interplay of the finite-time drive with QR—both processes known to yield nonequilibrium states—is found to produce an intriguing behavior of the defect density. The results reveal that QR increases the number of topological defects as the time scale τ increases, with a critical value at which the defect density abruptly changes its scaling with τ . We also find that the competition between adiabaticity and QR yields an optimal annealing time for which the defect density is at a minimum. This optimal time exhibits universal scaling with the QR rate, with the very same exponent as for scaling in the presence of noise in previous studies [16–18]. Hence, a system driven across, or terminating very close to a quantum critical point with Poissonian distributed QR remarkably appears to show the same anti-KZ scaling behavior as a ramp with uncorrelated Gaussian noise, indicating that universal mechanisms are in place for out-of-equilibrium dynamics near criticality.

Our results have been obtained by an exact solution for the time-dependent density matrix of the mode-decoupled two-band fermionic Hamiltonians of the problem, then adding QR to the linear ramp and numerically averaging over all QR realizations with a specified rate. This approach generalizes to any one-dimensional fermionic two-band model subject to

QR. We conjecture that crossovers from KZ to anti-KZ scaling induced by low-rate QR is generic for this full class of models, and possibly beyond.

The results presented hold promise as a testing ground for quantum annealing devices. Unlike environmental noise that also produce crossovers from KZ to anti-KZ scaling, QR protocols are expected to become programmable on future analog quantum computers, allowing for enhanced control. Such tests should be particularly relevant as it comes to the scaling of optimal annealing times with QR rates: Optimal times for quantum annealing identify the best compromise between maximal adiabaticity (slow ramp) and minimal decoherence (fast ramp)—a key figure of merit for adiabatic quantum computing.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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