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Recent theoretical and experimental works have explored universal dynamics related to surface growth physics in *isolated* quantum systems. In this Letter, we theoretically elucidate that *dissipation* drastically alters universal particle-number-fluctuation dynamics associated with surface-roughness growth in one-dimensional free fermions and bosons. In a system under dephasing that causes loss of spatial coherence, we numerically find that a universality class of surface-roughness dynamics changes from the ballistic class to a class with the Edwards-Wilkinson scaling exponents and an unconventional scaling function. We provide the analytical derivation of the diffusion equation from the dephasing Lindblad equation via a renormalization-group technique and succeed in explaining the drastic change. Furthermore, we numerically find the same change of the universality class under a more nontrivial dissipation, i.e., symmetric incoherent hopping.

Introduction. Universal dynamics has been explored for long years in classical statistical mechanics, and surface growth [1, 2] has been one of the most fundamental subjects for deepening our understanding of universal aspects behind the non-equilibrium phenomena. Minimal theoretical models for the classical surface growth are the Kardar-Parisi-Zhang (KPZ) [3] and Edwards-Wilkinson (EW) [4] equations, which exhibit universal dynamical scaling in the surface-height distribution [5, 6]. Recent works found a signature of KPZ universality even in isolated quantum many-body systems by investigating two-point spatio-temporal correlation functions numerically [7–14] and experimentally [15, 16]. Instead of computing the correlation function, a surface-height operator and quantum surface roughness were introduced in Refs. [17–19] by using the particle-number and spin fluctuations. Considering the particle-number fluctuations of isolated fermionic and bosonic lattice models, our previous works [17, 19] found emergence of the Family-Vicsek (FV) scaling [20, 21], the dynamical scaling of the surface roughness originally developed in classical surface growth [1]. As illustrated in Fig. 1(a), this scaling is characterized by three exponents α, β, and z, which determine universality classes of the dynamics [1].

In this Letter, we theoretically tackle a fundamental and intriguing question: "How does dissipation affect the universal fluctuation dynamics related to the surface-growth physics in quantum systems?" We consider open quantum systems with dissipation obeying the Lindblad equation (see Fig. 1(b)). To investigate large-scale long-time dynamics, we use an exact numerical method with correlation matrices [22–24] and a renormalization-group-based analytical technique [25–27]. First, studying free fermions and bosons on a one-dimensional (1D) lattice under dephasing, we numerically find that the FV scaling emerges even in the open quantum system and that the dissipation changes the universality class from the ballistic class to a class with the EW scaling

(a) Time evolution of surface roughness
(b) Schematics of our systems
(c) Table for the FV scaling exponents

Model	α	β	z
Edwards-Wilkinson model (EW class)	0.5	0.25	2
free fermions or bosons (Ballistic class)	0.5	0.5	1
free fermions + dephasing	0.490	0.245	2.00
free bosons + dephasing	0.498	0.264	1.91
free fermions + symmetric incoherent hopping	0.492	0.237	2.07
free bosons + symmetric incoherent hopping	0.508	0.263	1.94

this work

FIG. 1. (a) Schematic illustration for the Family-Vicsek (FV) scaling with the scaling exponents α, β, and z. The abscissa is time t, and the ordinate is the surface roughness $w(L, t)$, which is a function of a system size L and time t. The curves with the same color are the surface roughness with different system sizes. The roughness grows with t^β for $t \ll t^*$, where the saturation time t^* scales as L^z. The saturated surface roughness is proportional to L^α. (b) Schematic illustrations for (i) an isolated system and (ii) an open system considered in our work. (c) Table for the scaling exponents of the FV scaling. The first two rows are the previously known exponents, while the remaining four rows are the main results in this Letter. We here show the exponents for a staggered initial state.

exponents (see Fig. 1(c)) and an unconventional scaling function. Second, we analytically explain the change of the exponents by deriving diffusion equations via the renormalization-group method. To the best of our knowledge, this is the first analytical derivation of the diffusion equations from the dephasing Lindblad equation. Third, we use symmetric incoherent hopping as a non-

trivial dissipation, numerically finding the same change of the FV scaling exponents and the same unconventional scaling function. This strengthens the argument that the diffusive-type FV scaling is universal in open quantum systems under particle-number conservation. Figure 1(c) summarizes our results. Finally, we comment on absence of the FV scaling in Lindblad equations without the particle-number conservation and discuss experimental possibilities for observing our theoretical predictions.

Setup. We consider free fermions or bosons on a 1D lattice $\Lambda = \{1, 2, \cdots, L\}$ with an even number L of the lattice points. Let $\hat{a}_j$ and $\hat{a}_j^\dagger$ be the annihilation and creation operators at a site $j \in \Lambda$. When the particles are fermions (bosons), the operators satisfy $[\hat{a}_j, \hat{a}_k^\dagger]_+ = \delta_{jk}$ ($[\hat{a}_j, \hat{a}_k^\dagger]_- = \delta_{jk}$), where we introduce the (anti)commutator $[\hat{A}, \hat{B}]_\pm = \hat{A}\hat{B} \pm \hat{B}\hat{A}$ for operators $\hat{A}$ and $\hat{B}$. We assume that a quantum state at time t, specified by a density matrix $\hat{\rho}(t)$, obeys the Lindblad equation [28]:

$$\frac{d}{dt}\hat{\rho}(t) = -i[\hat{H}, \hat{\rho}(t)]_- + \mathcal{D}[\hat{\rho}(t)], \tag{1}$$

where $\hat{H}$ and $\mathcal{D}[\hat{\rho}(t)]$ are respectively a Hamiltonian and a dissipator. The Hamiltonian $\hat{H} = -\sum_{j=1}^{L-1} \left(\hat{a}_{j+1}^\dagger \hat{a}_j + \hat{a}_j^\dagger \hat{a}_{j+1} \right)$ describes coherent dynamics for the non-interacting particles. In this work, we consider two kinds of dissipators conserving the total particle number. One is the dephasing dissipator defined by

$$\mathcal{D}_{\mathrm{DEP}}[\hat{\rho}(t)] = \gamma \sum_{j=1}^{L} \left(\hat{n}_j \hat{\rho}(t) \hat{n}_j - \frac{1}{2}[\hat{n}_j^2, \hat{\rho}(t)]_+ \right) \tag{2}$$

with a particle-number operator $\hat{n}_j = \hat{a}_j^\dagger \hat{a}_j$ and a strength γ of the dephasing. The other is the dissipator for the symmetric incoherent hopping [23]:

$$\mathcal{D}_{\mathrm{SIH}}[\hat{\rho}(t)] = \gamma' \sum_{j=1}^{L} \left(\hat{R}_j \hat{\rho}(t) \hat{R}_j^\dagger - \frac{1}{2}[\hat{R}_j^\dagger \hat{R}_j, \hat{\rho}(t)]_+ \right)$$
$$+ \gamma' \sum_{j=1}^{L} \left(\hat{L}_j \hat{\rho}(t) \hat{L}_j^\dagger - \frac{1}{2}[\hat{L}_j^\dagger \hat{L}_j, \hat{\rho}(t)]_+ \right) \tag{3}$$

with the operators $\hat{R}_j = \hat{L}_j^\dagger = \hat{a}_{j+1}^\dagger \hat{a}_j$ and the parameter γ'. The first and second terms on the right hand side are responsible for the incoherent hopping from j to $j+1$ and vice versa, respectively. In this case, we use $\hat{H} = -\sum_{j=1}^{L} \left(\hat{a}_{j+1}^\dagger \hat{a}_j + \hat{a}_j^\dagger \hat{a}_{j+1} \right)$ under a periodic boundary condition.

The physical quantity of interest is surface roughness defined by a variance of a surface-height operator $\hat{h}_j = \sum_{k=1}^{j} (\hat{n}_k - \nu)$ with an initial filling factor ν [17–19]. This operator was introduced on the basis of a mathematical correspondence between a classical surface height $h(x,t)$ in the KPZ equation and a sound

mode $\delta n(x, t)$ in the fluctuating hydrodynamics in 1D systems. Since correlation functions for $\partial_x h(x, t)$ and $\delta n(x, t)$ have the same scaling function in the stationary processes, one can define an effective surface height by $h_{\mathrm{eff}}(x, t) = \int_0^x \delta n(y, t) dy$ in the fluctuating hydrodynamics. The surface-height operator $\hat{h}_j$ is quantum extension of $h_{\mathrm{eff}}(x, t)$. Using this surface-height operator, we define the surface roughness at a site $j \in \Lambda$ by $w_j(t) = \sqrt{\langle (\hat{h}_j - \langle \hat{h}_j \rangle_t)^2 \rangle_t}$ with the quantum statistical average $\langle \cdots \rangle_t = \mathrm{Tr}[\hat{\rho}(t) \cdots]$. In what follows, we focus on $j = L/2$ because the surface roughness grows for longer time for this choice than the other j and introduce the notation $w(L, t) = w_{L/2}(t)$ for brevity (see Sec. I of Supplemental Material (SM) [30]).

When the fluctuation of the surface height is scale-invariant, the surface roughness shows the FV scaling [1, 20, 21] defined by

$$w(L, t) = s^{-\alpha} w(sL, s^z t) \propto \begin{cases} t^\beta & (t \ll t^*) \\ L^\alpha & (t^* \ll t). \end{cases} \tag{4}$$

Here, the scaling exponents α, β, and z satisfying the scaling relation $z = \alpha/\beta$ classify a universality class, and t^* is saturation time. Two well-known universality classes originally found in classical systems are the KPZ [3] and EW classes [4] characterized by $(\alpha, \beta, z) = (1/2, 1/3, 3/2)$ and $(1/2, 1/4, 2)$, which show superdiffusive ($1 < z < 2$) and diffusive ($z = 2$) transport, respectively. In quantum systems, free fermions have $(\alpha, \beta, z) \simeq (1/2, 1/2, 1)$ [17, 19], which we call a ballistic class since the dynamical exponent z is unity.

Numerical method. Instead of directly solving the Lindblad equation, we solve the equations of motion for two- and four-point correlation matrices [22–24] defined by $D_{mn} = \langle \hat{a}_m^\dagger \hat{a}_n \rangle_t$ and $F_{mnpq} = \langle \hat{a}_m^\dagger \hat{a}_n^\dagger \hat{a}_p \hat{a}_q \rangle_t$, respectively. As shown in Sec. II of SM [30], we exactly derive the closed equations of motion, which enable us to access the long-time universal dynamics in the open quantum systems. Note that a third quantization and a superoperator method are other well-known efficient techniques to solve the Lindblad equation [31–38], but they are inconvenient in the dephasing case since this dissipation generates quartic terms in the thermofield representation.

Solving the equations for the correlation matrices, we can calculate the surface roughness using the following formula:

$$w(L, t)^2 = \pm \sum_{m=1}^{L/2} \sum_{n=1}^{L/2} F_{mnmn}(t) + (1 - \nu L) \sum_{m=1}^{L/2} D_{mm}(t)$$
$$+ \frac{\nu^2 L^2}{4} - \left(\sum_{m=1}^{L/2} D_{mm}(t) - \frac{\nu L}{2} \right)^2. \tag{5}$$

Here, $-(+)$ sign is for fermions (bosons).

FV scaling under the dephasing. We study how the dephasing (2) affects the surface-roughness dynamics. In this model, the dynamics occurs in a sector with a fixed
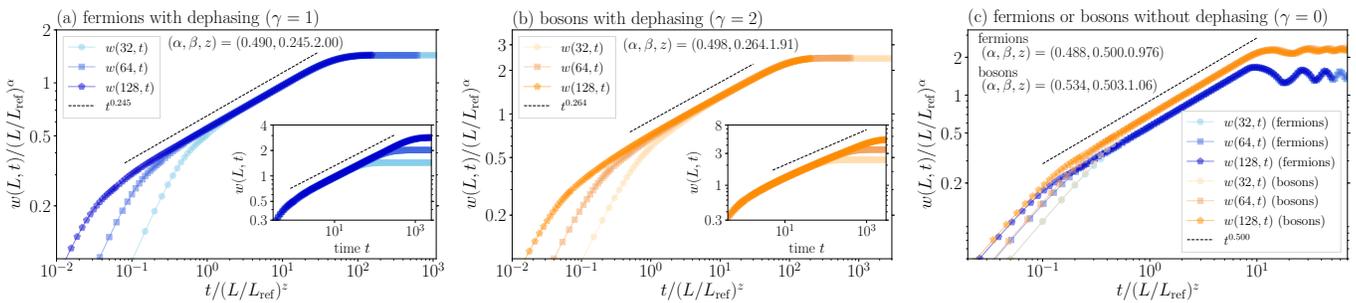
FIG. 2. Time evolution of the surface roughness for (a) fermions with $\gamma = 1$, (b) bosons with $\gamma = 2$, and (c) fermions and bosons with $\gamma = 0$. In the main panel, we show the surface roughness with the ordinate and abscissa normalized by $(L/L_{\text{ref}})^\alpha$ and $(L/L_{\text{ref}})^z$ with the reference system size $L_{\text{ref}} = 32$ [29]. The estimated exponents (α, β, z) are (a) $(0.490, 0.245, 2.00)$, (b) $(0.498, 0.264, 1.91)$, and (c) $(0.488, 0.500, 0.976)$ for the fermion and $(0.534, 0.503, 1.06)$ for the boson.

total particle number, and hence, we expect that the surface roughness comprised of the local particle number exhibits dynamics whose time scale increases with the system size L. The initial states used here are a staggered state (SS) $|\text{SS}\rangle = \prod_{j=1}^{L/2} \hat{a}_{2j}^\dagger |0\rangle$, a domain-wall state (DWS) $|\text{DWS}\rangle = \prod_{j=1}^{L/2} \hat{a}_j^\dagger |0\rangle$, and a uniform state (US) $|\text{US}\rangle = \prod_{j=1}^{L} \hat{a}_j^\dagger |0\rangle$ with the vacuum $|0\rangle$.

Figure 2 shows time evolution of the surface roughness for (a) fermions with $\gamma = 1$, (b) bosons with $\gamma = 2$, and (c) fermions and bosons with $\gamma = 0$. The initial state is SS. From Figs. 2(a) and (b), one can see that the FV scaling is well satisfied for both fermions and bosons with the dephasing, and the scaling exponents are almost the same as those for the EW class. Interestingly, we find that the scaling function has an unconventional form being different from the EW equation as discussed in Sec. III of SM [30] (see also Fig. 4(b)). On the other hand, the isolated fermions and bosons show the FV scaling with the ballistic class [17, 19] as shown in Fig. 2(c). Our numerical results clearly show that the dephasing alters the universality class from the ballistic class to the one with the EW-type exponents characterized by the diffusive dynamics with $z = 2$.

Next, we investigate dependence of the dynamics on γ. Figure 3(a) shows the time evolution of $w(L, t)$ for fermions with $\gamma = 1, 2^{-1}$, and 2^{-2}. We find that the surface roughness obeys $t^{0.25}$ in the late dynamics ($t \gtrsim 1/\gamma$), which corresponds to the EW exponent. This fact is clearly seen in the inset of Fig. 3(a). From this result, we argue that the change of the universality class occurs for infinitesimal dissipation strength γ, which indicates the strong impact of dissipation. To strengthen our argument, we also study whether the diffusive transport emerges in the dynamics starting from the DWS. As discussed in Ref. [16], the particle transfer from the left to right region is used to study the transport property. Here, we numerically compute $P_{\text{tra}}(t) = (N_R(t) - N_R(0)) - (N_L(t) - N_L(0))$ with $N_L(t) = \sum_{m=1}^{L/2} \langle \hat{a}_m^\dagger \hat{a}_m \rangle_t$ and $N_R(t) = \sum_{m=L/2+1}^{L} \langle \hat{a}_m^\dagger \hat{a}_m \rangle_t$. If the transport is dif-

fusive, we have $P_{\text{tra}}(t) \propto t^{0.5}$. Figure 3(b) shows the time evolution of $P_{\text{tra}}(t)$ for the fermions with $\gamma = 1, 2^{-2}, 2^{-4}$, and 2^{-6} in the large system size $L = 4096$ [39]. This result demonstrates that, for $\gamma \gtrsim 2^{-6}$, the ballistic behavior appears in the early dynamics ($1/\gamma \gtrsim t \gtrsim 1$) but the transport eventually becomes diffusive for sufficiently long time ($t \gtrsim 1/\gamma$). This numerical finding strongly supports our argument.

We numerically investigate the dependence of the FV scaling exponents on the initial states. The detailed time-evolution of the surface roughness is given in Sec. IV of SM [30]. The obtained exponents are summarized in Table I, which shows that the exponents are almost indepen-
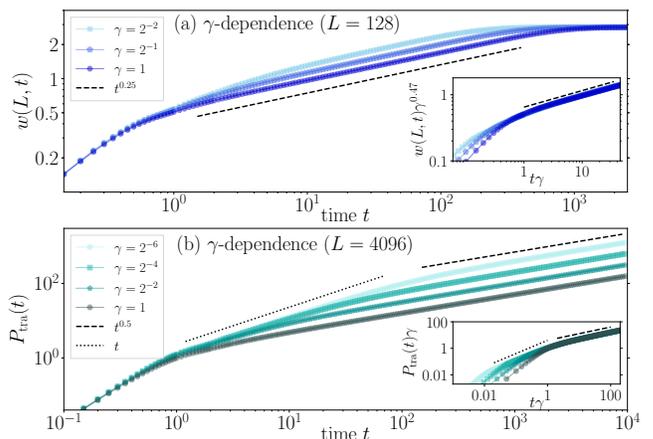
FIG. 3. Dependence of the fermionic dynamics on γ. (a) Time evolution of $w(128, t)$ for $\gamma = 1, 2^{-1}$, and 2^{-2}. The initial state is SS. In the inset, we normalize the time and the roughness by $1/\gamma$ and $1/\gamma^{0.47}$. This clearly exhibits that $\beta = 0.25$, the signature of the EW class (diffusive dynamics), emerges in $t \gtrsim 1/\gamma$. (b) Time evolution of $P_{\text{tra}}(t)$ for $\gamma = 1, 2^{-2}, 2^{-4}$, and 2^{-6}. The initial state is DWS, and the system size is 4096. In the inset, we normalize the time and the roughness by $1/\gamma$. Similar to (a), the diffusive behavior ($P_{\text{tra}}(t) \propto t^{0.5}$) emerges for $t \gtrsim 1/\gamma$.

dent of the initial states if the initial surface roughness is small.

Renormalization-group analysis. To understand the change of the universality class analytically, we use a perturbative renormalization-group method [25–27, 42] and derive effective equations for D_{mm} and F_{mnmn}, which determine the surface roughness (5). As derived in Sec. V of SM [30], when the dephasing strength γ is strong, the effective equations for the fermions and the bosons become

$$\frac{d}{dt}D_{mm} \simeq \frac{2}{\gamma}(D_{(m+1)(m+1)} + D_{(m-1)(m-1)} - 2D_{mm}), \quad (6)$$

$$\frac{d}{dt}F_{mnmn} \simeq \frac{2}{\gamma}(F_{(m+1)n(m+1)n} + F_{(m-1)n(m-1)n} + F_{m(n+1)m(n+1)} + F_{m(n-1)m(n-1)} - 4F_{mnmn})$$

$$(7)$$

for $|m - n| > 2$. Taking the continuum limit for these equations, we obtain the one- and two-dimensional diffusion equations, which are responsible for the diffusive transport. As far as we know, this is the first analytical derivation of the diffusion equations from the dephasing Lindblad equation. As discussed in Sec. V of SM [30], the effective equations show the similar dynamics to the exact numerical results. Our renormalization-group analysis can be applied to other Lindblad equations, offering an interesting method to study the Lindblad dynamics.

Our renormalization-group analysis explains the emergence of the EW exponents. The dynamical exponent z in the FV scaling is related to the correlation length $\xi(t) \propto t^{1/z}$ for the surface-height correlation [1]. Since the surface-height operator consists of the particle-number operator $\hat{n}_j$ and the effective equations support the diffusive particle transport, we expect $z = 2$. As discussed in Refs. [17, 19], we have $\alpha = 0.5$ in typical systems. Then, the scaling relation leads to $\beta = \alpha/z = 0.25$.

While the above method is only for strong γ, in Sec. VI of SM [30], making several assumptions based on the exponential decay of the off-diagonal elements of the correlation matrices and focusing on the late stage of the dy-

FIG. 4. Time evolution of the surface roughness under the symmetric incoherent hopping. Panel (a) displays the surface roughness for fermions ($\gamma' = 2$) and bosons ($\gamma' = 4$) obtained by the Lindblad equation with the symmetric incoherent hopping. The ordinate and abscissa are normalized by $(L/L_{ref})^{\alpha}$ and $(L/L_{ref})^z$ with $L_{ref} = 32$, and the estimated values of the FV scaling exponents are $(\alpha, \beta, z) = (0.492, 0.237, 2.07)$ for the fermions and $(0.508, 0.263, 1.94)$ for the bosons. Panel (b) shows the surface roughness for the four different Lindblad equations and the EW equation. The detail of the EW equation is given in Sec. III of SM [30]. The values of A and B for the Lindblad equations are determined such that all the curves in the late stages are collapsed to a single curve by eye. The result for the EW equation does not match the other curves for any values of A and B.

namics ($t \gtrsim 1/\gamma$), we derive the same effective equations without assuming the strong dephasing. This supports the emergence of the EW scaling exponents for infinitesimally small dephasing.

FV scaling under the symmetric incoherent hopping. We next show the emergence of the FV scaling in the Lindblad equation for the symmetric incoherent hopping (3). The numerical method is similar to that for the dephasing case, and we use the SS state as an initial state.

Figure 4(a) displays $w(L, t)$ for the fermions and the bosons. The FV scaling clearly appears even under the symmetric incoherent hopping. The estimated scaling exponents are very close to the EW class. Surprisingly, as shown in Fig. 4(b), the scaling functions are almost the same as those of the dephasing Lindblad equations irrespective of the particle statistics, and the form of the functions is different from the EW scaling function. Our findings strengthen the argument that the open quantum systems with the particle-number conservation universally exhibit the FV scaling with the EW scaling exponents but with the non-EW scaling functions.

Discussion. We have so far addressed the dissipation conserving the total particle number, but in general in-

TABLE I. FV scaling exponents for the dephasing model of fermions ($\gamma = 1$) and bosons ($\gamma = 2$) [40]. The obtained FV scaling exponents are close to the exponents of the EW class. The fitting error are 3σ error evaluated in the method of Ref. [41].

Initial state	α	β	z
[fermion]			
SS	0.490 ± 0.035	0.245 ± 0.005	2.00 ± 0.15
DWS	0.483 ± 0.061	0.249 ± 0.006	1.96 ± 0.26
[boson]			
SS	0.498 ± 0.067	0.264 ± 0.007	1.91 ± 0.29
DWS	0.503 ± 0.036	0.260 ± 0.005	1.95 ± 0.16
US	0.499 ± 0.052	0.262 ± 0.006	1.92 ± 0.23

teractions with environments lead to gain and loss of the particles. As shown in Sec. VII of SM [30], we numerically solve the Lindblad equation with the in and outflow of the particles at all the sites, finding absence of the FV scaling. We conjecture that the particle-number conservation can be essential for the emergence of the FV scaling because it ensures that the time scale of the surface roughness becomes larger as the system size increases. Another interesting case is a boundary-driven system, where the particles are injected (removed) at the left (right) edge, and thus the particles in the bulk are locally conserved. We will leave exploration of the FV scaling for this interesting setup.

We discuss experimental possibilities for observing our theoretical predictions. An experiment under photon scattering [43] is considered to be well described by the Lindblad equation for the dephasing (2). Thus, the change of universality class can be experimentally accessible when one observes dynamics beyond the dephasing timescale $1/\gamma$. As to the symmetric incoherent hopping, Refs. [44, 45] theoretically proposed how to realize it experimentally, and in future the universal scaling may be explored on the basis of these proposals.

Conclusion. We theoretically studied the surface-roughness dynamics described by the Lindblad equation with the two types of dissipation: one is the dephasing, and the other is the symmetric incoherent hopping. In both cases, we numerically found the emergence of a clear FV scaling. In the former case, we analytically elucidated that the universality class is altered due to the presence of the dephasing. From these numerical and analytical results, we argued that the change of the universality class occurs at infinitesimally small dissipation, suggesting the substantial impact of dissipation.

Our findings pave an intriguing avenue for exploring the universal fluctuation dynamics in open quantum systems. It is interesting to study whether open quantum many-body systems exhibit novel universal dynamics triggered by the interactions. It is also important to consider other dissipation such as asymmetric incoherent hopping [23, 46–50], which has close relation to classical stochastic processes, e.g., an asymmetric simple exclusion process [51, 52].

ACKNOWLEDGMENTS

[1] A.-L. Barabási and H. E. Stanley, *Fractal concepts in surface growth* (Cambridge university press, 1995).

[2] U. C. Täuber, *Critical dynamics: a field theory approach to equilibrium and non-equilibrium scaling behavior* (Cambridge University Press, 2014).

[3] M. Kardar, G. Parisi, and Y.-C. Zhang, Dynamic scaling of growing interfaces, Phys. Rev. Lett. **56**, 889 (1986).

[4] S. F. Edwards and D. R. Wilkinson, The surface statistics of a granular aggregate, Proceedings of the Royal Society of London. Series A, Mathematical and Physical Sciences **381**, 17 (1982).

[5] T. Kriecherbauer and J. Krug, A pedestrian's view on interacting particle systems, KPZ universality and random matrices, Journal of Physics A: Mathematical and Theoretical **43**, 403001 (2010).

[6] K. A. Takeuchi, An appetizer to modern developments on the kardarparisizhang universality class, Physica A: Statistical Mechanics and its Applications **504**, 77 (2018), lecture Notes of the 14th International Summer School on Fundamental Problems in Statistical Physics.

[7] M. Ljubotina, M. Žnidarič, and T. Prosen, Kardar-Parisi-Zhang physics in the quantum heisenberg magnet, Phys. Rev. Lett. **122**, 210602 (2019).

[8] J. De Nardis, M. Medenjak, C. Karrasch, and E. Ilievski, Anomalous spin diffusion in one-dimensional antiferromagnets, Phys. Rev. Lett. **123**, 186601 (2019).

[9] S. Gopalakrishnan, R. Vasseur, and B. Ware, Anomalous relaxation and the high-temperature structure factor of xxz spin chains, Proceedings of the National Academy of Sciences **116**, 16250 (2019).

[10] F. Weiner, P. Schmitteckert, S. Bera, and F. Evers, High-temperature spin dynamics in the heisenberg chain: Magnon propagation and emerging Kardar-Parisi-Zhang scaling in the zero-magnetization limit, Phys. Rev. B **101**, 045115 (2020).

[11] M. Dupont and J. E. Moore, Universal spin dynamics in infinite-temperature one-dimensional quantum magnets, Phys. Rev. B **101**, 121106 (2020).

[12] J. De Nardis, M. Medenjak, C. Karrasch, and E. Ilievski, Universality classes of spin transport in one-dimensional isotropic magnets: The onset of logarithmic anomalies, Phys. Rev. Lett. **124**, 210605 (2020).

[13] J. De Nardis, S. Gopalakrishnan, E. Ilievski, and R. Vasseur, Superdiffusion from emergent classical solitons in quantum spin chains, Phys. Rev. Lett. **125**, 070601 (2020).

[14] E. Ilievski, J. De Nardis, S. Gopalakrishnan, R. Vasseur, and B. Ware, Superuniversality of superdiffusion, Phys. Rev. X **11**, 031023 (2021).

[15] A. Scheie, N. E. Sherman, M. Dupont, S. E. Nagler, M. B. Stone, G. E. Granroth, J. E. Moore, and D. A. Tennant, Detection of Kardar–Parisi–Zhang hydrodynamics in a quantum heisenberg spin-1/2 chain, Nature Physics **17**, 726 (2021).

[16] D. Wei, A. Rubio-Abadal, B. Ye, F. Machado, J. Kemp, K. Srakaew, S. Hollerith, J. Rui, S. Gopalakrishnan, N. Y. Yao, I. Bloch, and J. Zeiher, Quantum gas microscopy of kardar-parisi-zhang superdiffusion, Science

376, 716 (2022).

[17] K. Fujimoto, R. Hamazaki, and Y. Kawaguchi, Family-Vicsek scaling of roughness growth in a strongly interacting bose gas, Phys. Rev. Lett. **124**, 210604 (2020).

[18] T. Jin, A. Krajenbrink, and D. Bernard, From stochastic spin chains to quantum Kardar-Parisi-Zhang dynamics, Phys. Rev. Lett. **125**, 040603 (2020).

[19] K. Fujimoto, R. Hamazaki, and Y. Kawaguchi, Dynamical scaling of surface roughness and entanglement entropy in disordered fermion models, Phys. Rev. Lett. **127**, 090601 (2021).

[20] T. Vicsek and F. Family, Dynamic scaling for aggregation of clusters, Phys. Rev. Lett. **52**, 1669 (1984).

[21] F. Family and T. Vicsek, Scaling of the active zone in the eden process on percolation networks and the ballistic deposition model, Journal of Physics A: Mathematical and General **18**, L75 (1985).

[22] V. Eisler, Crossover between ballistic and diffusive transport: the quantum exclusion process, Journal of Statistical Mechanics: Theory and Experiment **2011**, P06007 (2011).

[23] K. Temme, M. M. Wolf, and F. Verstraete, Stochastic exclusion processes versus coherent transport, New Journal of Physics **14**, 075004 (2012).

[24] D. Nigro, D. Rossini, and E. Vicari, Competing coherent and dissipative dynamics close to quantum criticality, Phys. Rev. A **100**, 052108 (2019).

[25] L. Y. Chen, N. Goldenfeld, and Y. Oono, Renormalization group theory for global asymptotic analysis, Phys. Rev. Lett. **73**, 1311 (1994).

[26] L.-Y. Chen, N. Goldenfeld, and Y. Oono, Renormalization group and singular perturbations: Multiple scales, boundary layers, and reductive perturbation theory, Phys. Rev. E **54**, 376 (1996).

[27] H. Chiba, Extension and unification of singular perturbation methods for odes based on the renormalization group method, SIAM Journal on Applied Dynamical Systems **8**, 1066 (2009).

[28] F. P. Heinz-Peter Breuer, *The Theory of Open Quantum Systems* (Oxford University Press, 2002).

[29] The reference system size is arbitrary, and we choose $L_{\mathrm{ref}} = 32$ to make the order of the ordinate in Fig. 2 unity.

[30] See Supplemental Material, which includes Refs. [53, 54], for (I) Spatial dependence of the surface roughness, (II) Numerical method, (III) Scaling function of the Edwards-Wilkinson equation, (IV) Numerical data for the growing surface roughness and its dependence on the initial states, (V) Renormalization-group derivation of the diffusion equations for the dephasing Lindblad equation, (VI) Derivation of the diffusion equations for the Lindblad equation without the strong-dephasing condition, and (VII) Surface-roughness dynamics with the in- and out-flow of the particles.

[31] T. Prosen, Third quantization: a general method to solve master equations for quadratic open fermi systems, New Journal of Physics **10**, 043026 (2008).

[32] P. Kos and T. Prosen, Time-dependent correlation functions in open quadratic fermionic systems, Journal of Statistical Mechanics: Theory and Experiment **2017**, 123103 (2017).

[33] M. V. Medvedyeva, F. H. L. Essler, and T. Prosen, Exact bethe ansatz spectrum of a tight-binding chain with dephasing noise, Phys. Rev. Lett. **117**, 137202 (2016).

[34] N. Shibata and H. Katsura, Dissipative spin chain as a non-hermitian kitaev ladder, Phys. Rev. B **99**, 174303 (2019).

[35] N. Shibata and H. Katsura, Dissipative quantum ising chain as a non-hermitian ashkin-teller model, Phys. Rev. B **99**, 224432 (2019).

[36] A. A. Dzhioev and D. S. Kosov, Super-fermion representation of quantum kinetic equations for the electron transport problem, The Journal of Chemical Physics **134**, 044121 (2011).

[37] A. A. Dzhioev and D. S. Kosov, Nonequilibrium perturbation theory in liouville–fock space for inelastic electron transport, Journal of Physics: Condensed Matter **24**, 225304 (2012).

[38] M. V. Medvedyeva and S. Kehrein, Power-law approach to steady state in open lattices of noninteracting electrons, Phys. Rev. B **90**, 205410 (2014).

[39] The quantity $P_{\mathrm{tra}}(t)$ can be computed only by the two-point correlation matrix $D_{mn}(t)$. Thus we can go to the large system size $L = 4096$.

[40] The exponents (α, z) and the exponent β are independently determined by the fitting.

[41] A. Piñeiro Orioli, K. Boguslavski, and J. Berges, Universal self-similar dynamics of relativistic and nonrelativistic field theories near nonthermal fixed points, Phys. Rev. D **92**, 025041 (2015).

[42] The renormalization-group technique used in this work is a kind of a singular perturbation method that has been developed in nonlinear science and mathematics.

[43] H. P. Lüschen, P. Bordia, S. S. Hodgman, M. Schreiber, S. Sarkar, A. J. Daley, M. H. Fischer, E. Altman, I. Bloch, and U. Schneider, Signatures of many-body localization in a controlled open quantum system, Phys. Rev. X **7**, 011034 (2017).

[44] S. Diehl, A. Micheli, A. Kantian, B. Kraus, H. P. Büchler, and P. Zoller, Quantum states and phases in driven open quantum systems with cold atoms, Nature Physics **4**, 878 (2008).

[45] T. Haga, M. Nakagawa, R. Hamazaki, and M. Ueda, Liouvillian skin effect: Slowing down of relaxation processes without gap closing, Phys. Rev. Lett. **127**, 070402 (2021).

[46] F. Carollo, J. P. Garrahan, and I. Lesanovsky, Current fluctuations in boundary-driven quantum spin chains, Phys. Rev. B **98**, 094301 (2018).

[47] T. Haga, M. Nakagawa, R. Hamazaki, and M. Ueda, Liouvillian skin effect: Slowing down of relaxation processes without gap closing, Phys. Rev. Lett. **127**, 070402 (2021).

[48] D. Bernard and T. Jin, Open quantum symmetric simple exclusion process, Phys. Rev. Lett. **123**, 080601 (2019).

[49] D. Bernard and L. Piroli, Entanglement distribution in the quantum symmetric simple exclusion process, Phys. Rev. E **104**, 014146 (2021).

[50] D. Bernard and T. Jin, Solution to the quantum symmetric simple exclusion process: The continuous case, Communications in Mathematical Physics **384**, 1141 (2021).

[51] F. Spitzer, Interaction of markov processes, Advances in Mathematics **5**, 246 (1970).

[52] T. M. Liggett, *Stochastic Interacting Systems: Contact, Voter and Exclusion Processes* (Springer, 1999).

[53] G. Barontini, R. Labouvie, F. Stubenrauch, A. Vogler, V. Guarrera, and H. Ott, Controlling the dynamics of an open many-body quantum system with localized dissipa-

tion, Phys. Rev. Lett. **110**, 035302 (2013).

[54] R. Labouvie, B. Santra, S. Heun, and H. Ott, Bistability in a driven-dissipative superfluid, Phys. Rev. Lett. **116**, 235302 (2016).