Editors' Suggestion

Unified Interface Model for Dissipative Transport of Bosons and Fermions

Y. Minoguchi,[1,*,†] J. Huber,[1,*] L. Garbe,[1,2,3,4] A. Gambassi,[5] and P. Rabl[1,2,3,4]

[1]*Vienna Center for Quantum Science and Technology, Atominstitut, TU Wien, 1020 Vienna, Austria*
[2]*Technische Universität München, TUM School of Natural Sciences, Physics Department, 85748 Garching, Germany*
[3]*Walther-Meißner-Institut, Bayerische Akademie der Wissenschaften, 85748 Garching, Germany*
[4]*Munich Center for Quantum Science and Technology (MCQST), 80799 Munich, Germany*
[5]*SISSA—International School for Advanced Studies and INFN, via Bonomea 265, 34136, Trieste, Italy*

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We study the directed transport of bosons along a one dimensional lattice in a dissipative setting, where the hopping is only facilitated by coupling to a Markovian reservoir. By combining numerical simulations with a field-theoretic analysis, we investigate the current fluctuations for this process and determine its asymptotic behavior. These findings demonstrate that dissipative bosonic transport belongs to the Kardar-Parisi-Zhang universality class and therefore, in spite of the drastic difference in the underlying particle statistics, it features the same coarse-grained behavior as the corresponding asymmetric simple exclusion process for fermions. However, crucial differences between the two processes emerge when focusing on the full counting statistics of current fluctuations. By mapping both models to the physics of fluctuating interfaces, we find that dissipative transport of bosons and fermions can be understood as surface growth and erosion processes, respectively. Within this unified description, both the similarities and discrepancies between the full counting statistics of the transport are reconciled. Beyond purely theoretical interest, these findings are relevant for experiments with cold atoms or long-lived quasiparticles in nanophotonic lattices, where such transport scenarios can be realized.

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Transport is one of the most common nonequilibrium processes and thus the subject of intense research in many areas of physics [1–7]. Of considerable interest in this context is the analysis of elementary paradigmatic models, which capture essential aspects of transport in terms of simplified abstractions of broader classes of physical processes. A prominent example in statistical mechanics is the asymmetric simple exclusion process (ASEP) [8–12], which describes random asymmetric hopping of hard-core particles or fermions on a one-dimensional lattice. This model has been the subject of intense research in statistical physics and many exact results led to a deeper understanding of the physics far from equilibrium [13–16]. Most important for this Letter, it has been shown that the ASEP can be mapped on a model of a randomly fluctuating interface [16] that is well described in terms of the celebrated Kardar-Parisi-Zhang (KPZ) equation [17]. This established an unexpected connection between these two seemingly unrelated fields and has generated a lot of additional interest in this topic recently [18–23].

In this Letter, we study the dissipative transport of bosonic particles along a one-dimensional lattice, which are described by an asymmetric simple *inclusion* process (ASIP) [24–28]. Physically, this model is motivated by recent experimental setups with cold atoms in optical lattices [31–35] and with long-lived photonic, polaritonic, and plasmonic excitations in nanophotonic structures [22,36–45], where such bosonic transport can be realized. In contrast to Pauli's exclusion principle, which motivates the ASEP for fermions, bosonic atoms or quasiparticles favor the hopping to neighboring sites that are already occupied. This property gives rise to an opposite, bosonically enhanced *inclusive* transport, which drastically changes the overall particle flow and can even lead to condensation phenomena [27,46,47] that are naturally forbidden for fermions. Surprisingly, the following analysis reveals that in spite of their fundamentally different exchange statistics, the transport of bosons belongs to the same KPZ universality class as transport of fermions. At a closer inspection, however, there are subtle but important differences between the two processes, which we resolve in terms a unified interface model.

Model—We consider a one dimensional (1D) lattice as shown in Fig. 1(a), where bosonic atoms or other quasiparticles hop incoherently between neighboring sites with asymmetric rates γ_+ and γ_-. This is the case, for example,

*These authors contributed equally to this work.
†Contact author: yuri.minoguchi@gmail.com

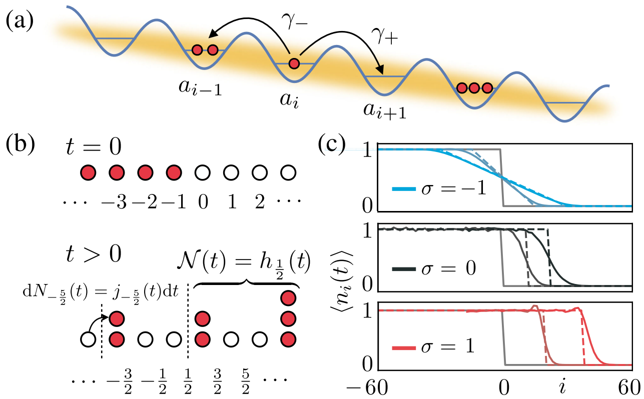

FIG. 1. (a) Sketch of the dissipative transport setup, where bosons hop along a tilted 1D lattice. Because of the coupling to a Markovian, low-temperature reservoir (yellow cloud), hopping becomes incoherent with asymmetric rates $\gamma_\pm$. (b) The lattice is initially prepared in a step initial condition and $\mathcal{N}(t)$ denotes the total number of particles at $i > 0$. The increments $\mathrm{d}N_{i+1/2} = j_{i+1/2}\mathrm{d}t$ count how many particles hop over the dual site $i + 1/2$ per unit time. (c) Comparison of the subsequent evolution of the average site occupation numbers $\langle n_i \rangle$ for fermions ($\sigma = -1$), bosons ($\sigma = +1$), and the ZRP ($\sigma = 0$), for $\gamma_+ = \gamma$ and $\gamma_- = 0$. The dashed line corresponds to the analytical solution of the (noiseless) Burgers equation.

for a reservoir-assisted tunneling processes in the presence of an external bias, and we refer to Refs. [27,48,49] for possible experimental realizations. We model the dynamics of this system by a quantum master equation for the density operator ρ,

$$\dot{\rho} = \gamma_+ \sum_{i=1}^{L-1} \mathcal{D}[a_{i+1}^\dagger a_i]\rho + \gamma_- \sum_{i=2}^{L} \mathcal{D}[a_{i-1}^\dagger a_i]\rho, \qquad (1)$$

where L is the total number of lattice sites, the a_i ($a_i^\dagger$) are bosonic or fermionic annihilation (creation) operators at site i and $\mathcal{D}[A]\rho = A\rho A^\dagger - (1/2)\{A^\dagger A, \rho\}$. For simplicity, we focus on the totally asymmetric case, $\gamma_+ = \gamma$ and $\gamma_- = 0$, but as long as $\gamma_+ \neq \gamma_-$, none of the following predictions crucially depend on this simplification.

In the absence of Hamiltonian dynamics, the density operator ρ evolving under Eq. (1) remains diagonal in the Fock basis $|\vec{n} = (n_1, \ldots, n_L)\rangle$, where $n_i = 0, 1, 2, \ldots$ is the number of bosons on site i. It is then sufficient to consider the dynamics of the diagonal elements $P_t(\vec{n}) = \langle \vec{n}|\rho(t)|\vec{n}\rangle$, which obey

$$\partial_t P_t(\vec{n}) = \gamma \sum_{i=1}^{L-1} [T_i^+ T_{i+1}^- - 1] n_i (1 + \sigma n_{i+1}) P_t(\vec{n}). \qquad (2)$$

Here, we introduced the operators $T_i^\pm f(\vec{n}) = f(\ldots, n_{i-1}, n_i \pm 1, n_{i+1}, \ldots)$ as well as the parameter $\sigma = 0, \pm 1$. The latter allows us to compare the bosonic ASIP ($\sigma = +1$) with the fermionic ASEP ($\sigma = -1$) and with a zero-range process (ZRP) ($\sigma = 0$) for independent classical particles, for which the hopping rate does not depend on the occupation of the target site [73,74].

The dynamics of the probability distribution $P_t(\vec{n})$ in Eq. (2) can be sampled by a Monte Carlo simulation of the site occupation numbers $n_i(t)$, which are updated at each time step $\mathrm{d}t$ as

$$n_i(t + \mathrm{d}t) = n_i(t) + \mathrm{d}N_{i-1/2}(t) - \mathrm{d}N_{i+1/2}(t). \qquad (3)$$

Here, $\mathrm{d}N_{i-1/2} = 0, 1$ are Poisson processes that indicate a particle hopping onto site i by jumping over the dual site $i - 1/2$, placed in between $i - 1$ and i [and similarly $\mathrm{d}N_{i+1/2}$ encodes the jumps between sites i and $i + 1$; see Fig. 1(b)]. The probabilities for these events to occur are $\Pr[\mathrm{d}N_{i+1/2}(t) = 1] = \Gamma_{i\to i+1}\mathrm{d}t$, where the hopping rates $\Gamma_{i\to i+1} = \gamma n_i (1 + \sigma n_{i+1})$ depend on the particle statistics σ. Note that Eq. (3) has the form of a continuity equation, from which the current $j_{i+1/2} = \mathrm{d}N_{i+1/2}/\mathrm{d}t$ can be directly read off.

Transport properties—Below we consider a *step initial condition* depicted in Fig. 1(b), where the left half of the lattice is filled with one particle per site and the right half is empty. By numerically evolving Eq. (3) and averaging over 10^3 realizations, we obtain the mean density profiles shown in Fig. 1(c). This plot shows that bosons (see lower plot) propagate into the empty region like a *shock wave*, which is accelerated compared to the diffusive advective transport of the ZRP (center plot) and the even slower dynamics of fermions (upper plot), which spread like a *rarefaction* wave (see [49] for details).

We now study the total number of transported particles $\mathcal{N}(t) = \sum_{i>0} n_i(t)$, which is shown for a few stochastic trajectories in Fig. 2(a). Overall, these trajectories increase linearly in time, with $\langle \mathcal{N}(t)\rangle \simeq j_\sigma t$ and an average current j_σ that depends on the particle statistics. The shot to shot fluctuations around this trend, $\Delta\mathcal{N}^2(t) = \langle \mathcal{N}^2(t)\rangle - \langle \mathcal{N}(t)\rangle^2$, are shown in Fig. 2(b) with a fit of the form $\Delta\mathcal{N}(t) \sim t^\beta$. For the ZRP the fluctuations increase with $\beta \simeq 1/4$, while a clearly distinct scaling exponent $\beta \simeq 1/3$ is observed for both fermions and bosons. Hence, in spite of their drastically different microscopic behavior reflected in the mean transport, these simulations indicate that bosons exhibit the same nontrivial scaling of fluctuations as fermions [14,75].

Fluctuating hydrodynamics—To explain these numerical findings, we consider the hydrodynamic limit $n_i(t) \to n(t, x)$, where we approximate the bosonic occupation numbers for the individual sites by a continuous density $n(t, x)$, coarse-grained over many adjacent lattice sites and time steps $\mathrm{d}t$. Applying either the Fokker-Planck approximation on Eq. (2) or the framework of nonlinear fluctuating hydrodynamics [76,77], we derive in Supplemental Material (SM) [49] an effective equation

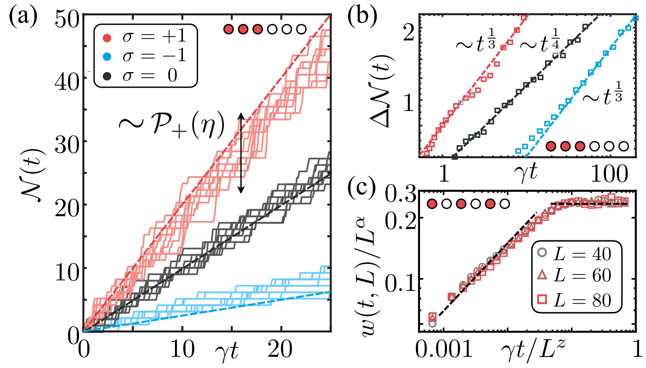

FIG. 2. (a) Individual realizations of the total number $\mathcal{N}(t)$ of transported particles, for different statistics σ and step initial conditions. The thick dashed lines indicate $\langle \mathcal{N}(t) \rangle = j_\sigma t$, as predicted by the noiseless Burgers equation. The corresponding currents are $j_+ = 2\gamma$ (bosons), $j_0 = \gamma$ (ZRP), and $j_- = \gamma/4$ (fermions). The fluctuations grow with t^β and so do the deviations from the noiseless Burgers equation. (b) Scaling of $\Delta\mathcal{N}(t)$ as a function of time. (c) Scaling collapse of $w(t, L)$ for ASIP for lattice size L, with periodic boundary conditions and initial condition with every other site occupied.

of motion for $n(t, x)$, which turns out to be the *stochastic Burgers equation* [78],

$$\partial_t n + \gamma(1 + 2\sigma n)\partial_x n = \nu \partial_x^2 n - \partial_x \xi. \tag{4}$$

We defined the viscosity $\nu = \gamma/2$ and a white noise $\xi \equiv \xi(t, x)$ with $\langle \xi(t, x)\xi(t', x')\rangle = D\delta(x - x')\delta(t - t')$ and $D = 2\nu n(x, t)[1 + \sigma n(x, t)]$. Since the precise magnitude of the noise does not alter the scaling of the current fluctuations, we set $D \simeq 2\nu\bar{n}(1 + \sigma\bar{n})$, where $\bar{n}$ is the steady state filling factor, in what follows. For $\sigma = -1$, Eq. (4) is then a well established effective description of the ASEP [75,79,80]. By omitting the viscosity, we obtain a deterministic linear transport equation for the ZRP ($\sigma = 0$) with current $j_0 = \gamma$. In case of bosons (fermions) we obtain the noiseless Burgers equation, which predicts a shock wave [81] (rarefaction wave [11,82]) and a current $j_+ = 2\gamma$ ($j_- = \gamma/4$) that matches the density profile in Fig. 1(c) [see dashed line in Fig. 2(a) and [49] for more details].

KPZ universality—Instead of characterizing the particle dynamics via the local site occupations n_i, we follow a well-known procedure [11,18–23,83] and consider the number of particles transported across the dual lattice site, $h_{i+1/2}(t) = \int_0^t \mathrm{d}\tau j_{i+1/2}(\tau)$ [see Fig. 1(b)]. The number of particles on site i is then the difference of particles having hopped on and off site i up to time t, i.e., $n_i(t) = h_{i-1/2}(t) - h_{i+1/2}(t)$. In this representation we identify $\mathcal{N}(t) \equiv h_{1/2}(t)$ and, in the hydrodynamic limit, $h_{i+1/2}(t) \to h(t, x)$, we obtain the continuity equation $-\partial_x h(t, x) = n(t, x)$. Inserting this relation into Eq. (4) yields

$$\partial_t h = v_\sigma - c_\sigma \partial_x h + \gamma\sigma(\partial_x h)^2 + \nu \partial_x^2 h + \xi, \tag{5}$$

with $c_\sigma = \gamma$ and $v_\sigma = 0$, which we introduced for later convenience. Up to a Galilean transformation, this is the well-known KPZ equation [17,79]. Note that the particle statistics σ determines the sign of the KPZ nonlinearity. The physics of this equation is characterized by the scaling exponent z, which governs the temporal growth of the correlation length $L_{\mathrm{corr}} \sim t^{1/z}$ and α characterizes the spatial roughness of the local transport fluctuations. Scaling theory then suggests that $\Delta h(t, 0) = \Delta\mathcal{N}(t) \sim t^{\alpha/z}$, or $\beta = \alpha/z$ in the notation above. In 1D the scaling exponents $z = 3/2$ and $\alpha = 1/2$ are known and independent of the sign σ of the nonlinearity (see [49] for details). This explains the identical exponents $\beta = 1/3$ obtained in numerical simulations for both fermions and bosons [51]. In contrast, for the Edwards-Wilkinson equation [84] obtained for the ZRP ($\sigma = 0$), α remains the same, but the dynamical scaling changes to $z = 2$, consistent with the numerically extracted exponent $\beta = 1/4$. To investigate scaling with the system size, we perform additional numerical simulations of the ASIP on a periodic chain, with translationally invariant initial conditions, i.e., with an alternating filling from site to site. In Fig. 2(c) a scaling collapse of $w(t, L) = \langle L^{-1}\sum_i [h_i(t) - \bar{h}(t)]^2\rangle^{1/2}$, with $\bar{h}(t) = L^{-1}\sum_i h_i(t)$, is shown for the theoretically predicted values of α and z, with Family-Viczek scaling $w(t, L) = L^\alpha f(t/L^z)$ [85,86]. This plot shows that both bosonic and fermionic transport obey the same KPZ scaling, and are thus much more similar than expected from the dynamics of averaged quantities.

Full counting statistics—Let us now go beyond the scaling of second moments and take a closer look at the statistics of $\mathcal{N}(t)$, which at long times is written in the asymptotic form [14]

$$\mathcal{N}(t) = j_\sigma t + (\Gamma_\sigma t)^\beta \eta. \tag{6}$$

Here, the first term describes the average linear growth with the current j_σ predicted by the noiseless Burgers equation. The second term describes the fluctuations with a scaling factor $\Gamma_\sigma = \gamma[\bar{n}_\sigma(1 + \sigma\bar{n}_\sigma)]^2 C_\sigma$, where $\bar{n}_\sigma$ corresponds to the long-time local density at the origin [see Fig. 1(c)]. For fermions with $\bar{n}_- = 1/2$ we find $C_- = 1$, which agrees with the exact solution [14], while for bosons with $\bar{n}_+ = 1$, our simulations suggest $C_+ \simeq 0.3$ (see [49]). The random variable η is distributed according to the *time-independent* probability distribution $\mathcal{P}_\sigma(\eta)$, which determines the full counting statistics (FCS) of $\mathcal{N}(t)$.

Starting from the step initial conditions, the distribution of $\mathcal{N}(t)$ can be extracted numerically for arbitrary σ. For the ZRP, the fluctuations are Gaussian due to the linearity of the equation. For the case of fermions, $\sigma = -1$, we find the *sign-reversed* Tracy-Widom distribution (TWD) [87] of the Gaussian unitary ensemble (GUE) [53,88] i.e., $\mathcal{P}_-(\eta) = \mathrm{TW}_{\mathrm{GUE}}(-\eta)$ [see Fig. 3(a)] in agreement with the exact solution [14]. Surprisingly, for the case of bosons, $\sigma = +1$,

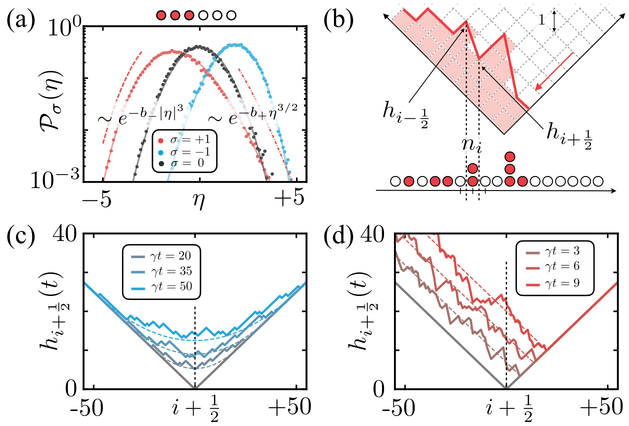

FIG. 3. (a) The numerically obtained FCS (dots) for the bosons (red), fermions (blue), and ZRP (black) for a step initial condition. The solid lines corresponds to $\mathrm{TW}_{\mathrm{GOE}}(\eta)$ (red), $\mathrm{TW}_{\mathrm{GUE}}(-\eta)$ (blue), and Gaussian (black). (b) Mapping the bosonic transport to corner growth. The red arrow corresponds to the direction of particle deposition (red squares) leading to a growth of the surface height $h_{i+1/2}$ (thick red line). Particle transport (solid lines) for fermions (c) and bosons (d) is plotted in terms of $h_i(t)$ for a single realization. The colored dashed lines correspond to the noise and noiseless solution of Eq. (5). The height along the black dashed line corresponds to $\mathcal{N}(t)$. In the fermionic case (c), the surface develops a curvature, whereas in the bosonic case (d) it stays flat.

and with the *same* initial conditions, we find the TWD of the Gaussian orthogonal ensemble (GOE) [89] *without* sign reversal i.e., $\mathcal{P}_+(\eta) = \mathrm{TW}_{\mathrm{GOE}}(\eta)$ [see Fig. 3(a)]. Accordingly, the FCS of the bosons is sign-reversed and is slightly broader (GOE) compared to its fermionic counterpart (GUE). Note that the negative (positive) first moment of the GOE (GUE) TWD implies that bosons (fermions) flow slightly slower (faster) than predicted by noiseless hydrodynamics [see Fig. 2(a)]. For periodic and stationary initial conditions the FCS coincide up to sign reversal [49].

Large-deviations analysis for bosons—It is possible to obtain a physical intuition for the emergence of asymmetric tails in the FCS of the bosons [90]. Let us consider rare events, where up to time t either no particle, $\mathcal{N}(t) \sim 0$, or an excess number of particles, $\mathcal{N}(t) \sim j_{\mathrm{exc}} t$ with $j_{\mathrm{exc}} \gg j_+$, were transported. In both cases we find from Eq. (6) the scaling $\eta \sim t^{2/3}$.

With this in mind we start our discussion by considering the right tail $\mathcal{P}_+(\eta \to \infty)$, i.e., the extreme case of an unusually large number of particles being transported. We assume that these events are dominated by shock configurations with a large number of particles $\bar{n}_{\mathrm{s}} \gg 1$ concentrated in the front. This gives rise to an excess current $j_{\mathrm{exc}} = \gamma \bar{n}_{\mathrm{s}}(1 + \bar{n}_{\mathrm{s}})$. The probability that in a single time step Δt, on hydrodynamics scales, a number of $j_{\mathrm{exc}}\Delta t$ particles is transported to the right is p_1. Since the shape of the shock front does not change, this probability remains constant and assuming independence of this excess transport in every time step we obtain $\Pr[\mathcal{N}(t) = j_{\mathrm{exc}} t] \sim (p_1)^{t/\Delta t} \sim \exp(-a_+ t)$, where the unknown constant a_+ depends on microscopic details. Going from t to the variable η we obtain the asymptotic form $P_+(\eta \to \infty) \sim \exp(-b_+ \eta^{3/2})$ [see red dashed dotted line in Fig. 3(a)].

The left tail $\mathcal{P}_-(\eta \to -\infty)$ is instead obtained by focusing on rare events where no particle is transported across the origin. Note that due to the ongoing transport up to the origin and their bosonic nature, particles pile up like $n_0(t) \sim j_+ t$ at the origin, and the probability of not jumping further right is progressively suppressed with the occupation $p(n_0) \simeq 1 - n_0 \gamma \Delta t \simeq (p_0)^{n_0}$ with $p_0 = 1 - \gamma \Delta t$. Then the probability that for long times t no particle hops to the right can be estimated as $\Pr[\mathcal{N}(t) = 0] \sim p_0^1 p_0^2 \cdot \ldots \cdot p_0^{t/\Delta t} \sim \exp(\sum_{i=1}^{t/\Delta t} i) \sim \exp(-a_- t^2)$. Going again from t to η we obtain $\mathcal{P}_+(\eta \to \infty) \sim \exp(-b_- \eta^3)$ [see red dashed line in Fig. 3(a)]. While these arguments are too heuristic to be able to predict $b_\pm$, they capture the correct scaling of the FCS tails, and show that bosonic statistics has a decisive effect on the shape (sign reversal) of the distribution.

Unified interface model—As a final step, we show how the differences and similarities between bosonic and fermionic transport can be reconciled by mapping both processes onto a common model of a fluctuating interface. Above, we showed how the transport process can be described in terms of a field $h_{i+1/2}$, which increases by 1 if a particle jumps from i to $i+1$. The key step is to interpret this quantity as the *height* of a randomly increasing interface. To make connection with the exact results for the ASEP [14], we exploit the freedom in the initial conditions of the height field and set $h_{i-1/2}(t) - h_{i+1/2}(t) = n_i(t) - 1/2$. With this convention, every particle configuration can be mapped to a height configuration of this effective interface model [see Fig. 3(b) for bosons and [14,16] for fermions] with an initial profile in the shape of a *corner*, $h_{i+1/2}(t=0) = |i|/2$. The number of particles transported to the right corresponds to the interface height at the origin, $\mathcal{N}(t) = h_{1/2}(t)$.

In the hydrodynamic limit the dynamics of the interface height are described by the KPZ equation (5) with a *deposition rate* $v_\sigma = (\gamma/2)(1 + \sigma/2)$ and advection $c_\sigma = \gamma(1 + \sigma)$. Formally, we obtain this by setting $h \to h + x/2$ in Eq. (5) [49]. In the case of bosons, there is an additional advection term c_+, which is absent for fermions. This term is responsible for the deposition of particles under an angle parallel to the right wedge [see red arrow in Fig. 3(b)] leading to an accumulation of particles on the left slope only. Finally, the positive (negative) sign of the KPZ nonlinearity translates into a *surface growth* (*surface erosion*) process. When combined, these three mechanisms result in two very distinct growth patterns, which are depicted in Figs. 3(c) and 3(d) for fermions and bosons, respectively.

Based on this interface model, the differences in the FCS of bosonic and fermionic transport can now be understood as follows. First, neglecting v_σ and c_σ for now, the transformation $h \to -h$ flips the sign of the KPZ nonlinearity, from growth to erosion, $\gamma \to -\gamma$, or maps bosons ($\sigma = 1$) into fermions ($\sigma = -1$). The transmutation of bosons into fermions under this transformation explains why the FCS are, up to a broadening, sign flipped with respect to each other [see Fig. 3(a)]. Secondly, we invoke the *refined* classification of the KPZ universality to explain the relative broadening of the FCS between bosons and fermions. This refined classification, which was originally conjectured by Prähofer and Spohn [11,91] and confirmed both theoretically and experimentally for many different examples, states that the KPZ universality class decomposes into subclasses depending on the geometry of the interface: the FCS of interfaces that are curved (flat) are of the GUE [14,91–95] (GOE [91,95–97]) type, respectively (see [49] for details). As discussed earlier, while the sign of the KPZ nonlinearity does not affect the scaling exponents, our analysis shows that the curvature is very much affected. This can already be anticipated from the KPZ equation in Eq. (5) in the absence of noise [dashed lines in Figs. 3(c) and 3(d)]. For the case of surface growth (erosion) the effective interface corresponding to bosonic (fermionic) transport stays asymptotically flat (curved) around $x = 0$. Therefore, provided that the Prähofer-Spohn conjecture holds, this interface model explains the different type of distributions observed in $\mathcal{N}(t) = h(t, x = 0)$. This in combination with the scaling analysis above and the additional simulations in [49] demonstrates that the ASIP is consistent with KPZ universality and all its subclasses.

Conclusion—In summary, we have analyzed the ASIP describing the directed dissipative transport of bosonic particles along a 1D lattice. Despite their drastically different hydrodynamics behavior, we have found surprising similarities between fermionic and bosonic transport in the scaling of fluctuations. However, we have also shown that these two cases map onto two different interface processes, a curved, eroding surface for fermions on the one hand and a flat, growing surface for bosons on the other hand. This mapping explains the subtle difference in the full counting statistics of the transport in full consistency with the KPZ universality class. These fundamental properties, which are exclusively attributed to the particle statistics, can be directly probed, i.e., in cold-atom experiments, where techniques to track the dynamics of individual particles are already available.

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