

Nonequilibrium Phases and Quantum Correlations in Synthetic Transport Models

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Quantum devices featuring midcircuit measurement and reset capabilities, such as quantum computers and dual-species Rydberg quantum simulators, enable the realization of *quantum cellular automata*. These systems evolve in discrete time following local updates implemented by unitary gates and allow for the realization of both closed and synthetic dissipative dynamics. Here, we focus on quantum cellular automata that implement minimal models of classical and quantum transport. To illustrate our ideas, we focus on a discrete-time totally asymmetric simple exclusion process and investigate how coherent dynamical contributions allow for the emergence of quantum effects and correlations. We find that bipartite entanglement dominates the transient evolution, while stationary states can retain quantum correlations beyond entanglement. Our results suggest viable routes for realizing transport models on quantum devices and characterizing collective quantum correlations in strongly driven systems.

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Introduction—Transport phenomena provide a unifying framework for the study of nonequilibrium systems. At the microscopic scale, particle transport is often modeled using *simple exclusion processes*, which describe the stochastic dynamics of mobile agents on lattices and complex networks. These processes have found a wide range of applications including RNA polymerization [1–5], transport through networks [6,7], and traffic modeling [8,9]. Among the many variants, the asymmetric simple exclusion process (ASEP) and its totally asymmetric counterpart (TASEP) have emerged as paradigmatic nonequilibrium models, characterized by a rich phenomenology, multiple stationary phases, with the stationary state exactly solvable in the classical case [10–28].

Cellular automata (CA) are minimal discrete-time models whose fundamental units, so-called cells, are updated according to local rules depending on the state of the neighboring cells [29,30]. Famous examples of CA include Conway’s Game of life [31] or Wolfram’s Rule 30 automaton [29,32], both of which are of particular interest as they produce complex behavior from seemingly simplistic rules. Their ability to generate emergent complexity motivates the study of CA dynamics and makes them a versatile tool for designing synthetic dynamical systems.

Because of their locality, CA provide a natural minimal framework for modeling simple exclusion processes.

CA have also been extended to the quantum setting. Here, each cell consists of a qubit (or, more generally, a d -level quantum system, known as a qudit), and the update rules are replaced by local unitary gates [33]; see Fig. 1(a).

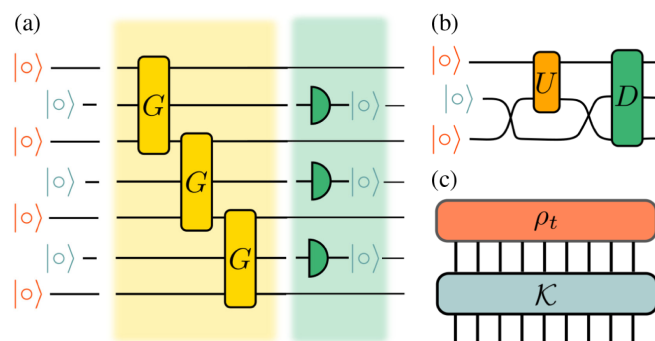


FIG. 1. Quantum cellular automaton from dual-species atom arrays. (a) Two-species atom array (sketched on the left) modeled of system qubits (leftmost in red) and ancillary qubits (rightmost in green). System and ancillary qubits interact through the repeated application of the local quantum gate G . After a full sweep throughout the lattice, the ancilla qubits are measured and reinitialized, and the sweep is repeated. (b) The gate G is composed of two unitary gates: U encodes coherent interactions within the system; D induces ancilla-mitigated dissipation in the system. SWAP operations are inserted to stress the nearest-neighbor range of the gate U and to ensure the final qubit order is consistent with the initial qubit order; they do not affect the effective transport dynamics. (c) The full QCA update, consisting of sweep and ancilla reinitialization, implements a dissipative dynamics that can be expressed as a global Kraus map $\mathcal{K}$ that propagates the quantum state through time $\rho_{t+1} = \mathcal{K}[\rho_t]$.

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Quantum CA (QCA) have recently been realized experimentally in programmable atomic arrays [34]. Midcircuit measurements and qubit resets enable the implementation of incoherent stochastic processes, allowing one to engineer synthetic open-system dynamics. In these settings, QCA have been identified as a potential approach to overcoming challenges associated with reliable manipulation of large many-qubit systems, as their simple structure can limit the complexity of the coherent dynamical control [34,35]. Furthermore, QCA have been utilized, both theoretically and experimentally, to explore nonequilibrium dynamics [36–38] as well as phase transitions into absorbing states [39] and collective properties in quantum neural networks [40–43]. However, QCA have not been used for exploring quantum transport, which is the main purpose of this Letter, particularly since discrete-time dynamics can lead to qualitatively different correlation structures [44].

Recently, quantum versions of simple exclusion processes have attracted significant interest [45–50]. Notably, the TASEP has emerged as a fundamental building block of *active quantum matter* [51]. Despite this progress, the role of genuine quantum correlations in driven exclusion dynamics remains poorly understood, particularly in the nonequilibrium steady states (NESS) arising from the interplay between coherent and stochastic transport.

In this Letter, we introduce a QCA realization of the TASEP and explore the competition of classical stochastic and coherent quantum transport. Recent experiments in programmable Rydberg atom arrays have demonstrated scalable control of many-body quantum dynamics, placing the systems considered here within experimental reach [52–55]. We investigate the nonequilibrium phase diagram of this QCA implementation of the quantum TASEP through large-scale simulations and characterize the resulting NESS. Strikingly, although standard entanglement witnesses do not detect quantum correlations, the steady states exhibit clear quantum signatures beyond entanglement. Given the simplicity and paradigmatic nature of the TASEP, our results indicate that quantum features beyond entanglement can appear even in minimal models of dissipative transport.

Transport QCA model—We consider the quantum circuit model depicted in Fig. 1(a). Each lattice site hosts a qubit, whose Z eigenstates are interpreted as occupied $|\bullet\rangle$ or vacant $|\circ\rangle$, representing agents in a one-dimensional lattice. Two neighboring qubits are paired with one ancillary qubit to form a *cluster*. The automaton rule, realized by the quantum gate G , is applied on these clusters. We explicitly consider a gate that consists of the subsequent application of a two-qubit gate U followed by a three-qubit gate D ; see Fig. 1(b). The gate U acts solely on the system qubits, whereas D allows the ancilla to interact with the system qubits. As depicted in Fig. 1(a), all ancilla qubits are measured and reinitialized once the automaton rule has been applied throughout the lattice, thereby inducing

dissipative processes in the system. For the clusters, we introduce the eight quantum states $|\mathfrak{g}\rangle$ where the ancilla label is chosen to sit on top of the two system labels, respecting the triangular structure introduced in Fig. 1. Upon measuring the ancilla, the QCA rule implements a dissipative quantum channel in each system cluster which is described by two Kraus operators $K_{\mathfrak{g}} = \langle \mathfrak{g} | D | \circ \rangle$ [56] corresponding to the measurement outcome. This setup allows us to include a single dissipative processes per cluster in the otherwise coherent QCA. An ancilla flip indicates if the corresponding dissipative process has occurred in a measured realization. This scheme can also be thought of as a repeated interaction framework [60], in which successive layers of ancillary qubits interact with a system through the repeated application of the QCA rule, implemented by a quantum gate [37,41,61–63].

To explore quantum transport phenomena in this QCA framework, we choose the automaton rule G as follows. The two-qubit system gate is the parametrized gate [64]

$$U(\omega) = |\circ\circ\rangle\langle\circ\circ| + (\cos\omega|\circ\bullet\rangle - i\sin\omega|\bullet\circ\rangle)\langle\circ\bullet| + (\cos\omega|\bullet\circ\rangle - i\sin\omega|\circ\bullet\rangle)\langle\bullet\circ| + |\bullet\bullet\rangle\langle\bullet\bullet|, \quad (1)$$

which implements coherent nearest-neighbor hopping of excitations. The system-ancilla component is

$$D(\tau) = \mathbb{1} + i\sqrt{\tau}\left(|\overset{\circ}{\bullet}\overset{\circ}{\circ}\rangle\langle\overset{\circ}{\bullet}\overset{\circ}{\bullet}| + |\overset{\circ}{\bullet}\overset{\circ}{\bullet}\rangle\langle\overset{\circ}{\circ}\overset{\circ}{\bullet}|\right) - (1 - \sqrt{1 - \tau})\left(|\overset{\circ}{\bullet}\overset{\circ}{\bullet}\rangle\langle\overset{\circ}{\bullet}\overset{\circ}{\bullet}| + |\overset{\circ}{\circ}\overset{\circ}{\bullet}\rangle\langle\overset{\circ}{\bullet}\overset{\circ}{\bullet}|\right), \quad (2)$$

which implements the stochastic TASEP update rule upon tracing over the ancilla [56]. Since the ancilla is in the vacant state at the start of each time update, we see that, apart from the identity, only two terms in $D(\tau)$ contribute. These correspond to a hopping transition term if the ancilla flips with probability τ $\left(|\overset{\circ}{\bullet}\overset{\circ}{\bullet}\rangle\langle\overset{\circ}{\circ}\overset{\circ}{\bullet}|\right)$ and a no-hop term if the ancilla does not flip $\left(|\overset{\circ}{\circ}\overset{\circ}{\bullet}\rangle\langle\overset{\circ}{\bullet}\overset{\circ}{\bullet}|\right)$. Hence, in the

absence of the gate $U(\omega)$, the repeated ancilla interaction implements the bulk rules of a discrete-time TASEP with the bulk hopping probability τ . To realize open-boundary TASEP dynamics, we consider a boundary driven system [57,65–69]. At the left boundary, excitations are injected with probability α , while at the right boundary they are ejected with probability β . These boundary processes can be implemented using two additional ancillary qubits that interact locally with the left and right boundary qubit, respectively, in direct analogy to the ancilla-mediated bulk hopping; see Supplemental Material for further details [56].

The stationary state for the classical TASEP can be analytically derived using a matrix product ansatz [15,56] and supports three distinct nonequilibrium phases. A

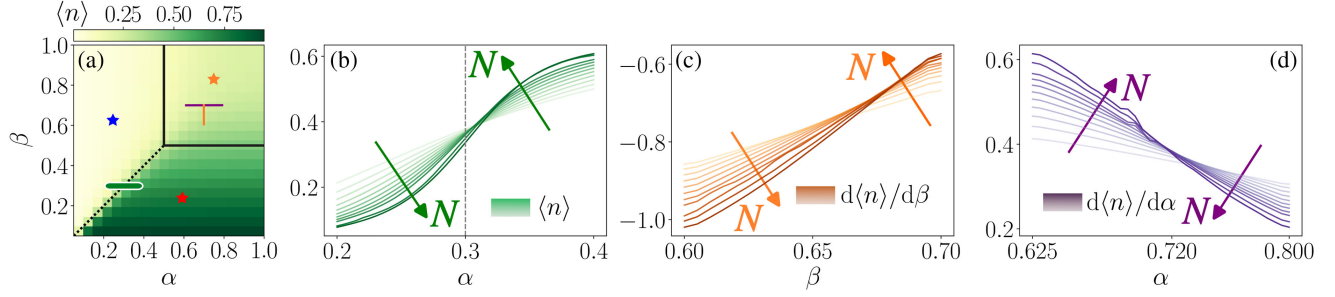


FIG. 2. Average lattice occupation in finite-size NESS. (a) Phase diagram for a system of $N = 30$ lattice sites with coherent transport induced by the gate $U(\pi/4)$ and bulk hopping probability $\tau = 0.75$. The dashed line indicates the classical coexistence line separating the LD and HD phases, and the solid line indicates the position of the transition to the MC phase in the classical model ($\omega = 0$). Colored markers indicate parameter values we further study in Fig. 3, and colored lines indicate parameter ranges for the order parameter scalings in (b)–(d). (b) shows the finite-size scaling of the average lattice occupation across the coexistence line ($\beta = 0.3$), and (c) and (d) show the finite-size scaling of the derivative of the average lattice occupation across the LD-MC ($\beta = 0.7$) and the HD-MC transition ($\alpha = 0.7$), respectively. All panels depict lattice sizes $N = 6, 8, 10, 12, 14, 16, 18, 20, 22, 26, 30$ encoded by the opacity of the lines with larger N being more opaque.

high-density phase (HD) is separated from a low-density phase (LD) by the coexistence line $\alpha = \beta$ up to $\alpha^* = \beta^* = 1 - \sqrt{1 - \tau}$. For $\alpha, \beta > \alpha^*, \beta^*$, the TASEP supports a maximum current phase (MC). While both the LD-MC and the HD-MC transition are continuous, the LD-HD transition is discontinuous [15,56]. For the remainder of this Letter, we choose $\tau = 0.75$.

Quantum NESS phases—We explore the influence of coherent transport induced by $U(\omega)$ on the phase diagram of the classical TASEP. We note that $U(\omega)$ is the identity in the zero- and the two-particle sector of each cluster and acts nontrivially only in the single-particle sector. For $\omega = \pi/4$, this gate can create maximally entangled Bell states between neighboring system sites, such that the influence of the coherent transport is expected to be most significant. Thus, we restrict our analysis to $\omega = \pi/4$ for the remainder of this Letter.

We simulate the QCA model using a time-evolving block decimation tensor-network algorithm [70,71], allowing the bond dimension to adapt dynamically during the evolution. A single time step consists of the iterative application of the left boundary drive, the sequential bulk update, and the right boundary drive; see Supplemental Material for additional details [56]. This procedure allows us to time evolve the system to its NESS determined by the boundary driving parameters α and β . Figure 2 shows our findings for systems up to $N = 30$ qubits.

We see from Fig. 2(a) that the general shape of the classical phase diagram stays intact upon the introduction of strong coherent transport. In Figs. 2(b)–2(d), we illustrate the relevant order parameter crossings of the average lattice density that define the exact location of the phase boundaries. Upon increasing the system size N across the LD-HD coexistence line, the average occupation decreases for smaller α and increases for larger α . This behavior is indicated by a crossing of the finite-size curves and suggests the emergence of a jump discontinuity; see

Fig. 2(b). Conversely, the LD-MC as well as the HD-MC transition remain continuous, and only the derivative of the average number density develops a discontinuity. However, the extent of the MC phase is diminished compared to the classical case as can be seen from the finite-size scaling analysis in Figs. 2(c) and 2(d). As we explained above, the onset of the MC phase in the classical case depends on the numerical value of the bulk hopping probability τ . Larger τ values diminish the size of the MC phase, similar to what we observe in the presence of coherent transport. Thus, the coherent transport appears to renormalize the bulk hopping probability toward higher values, which we further analyse in the Supplemental Material [56]. The collective behavior of the quantum model seems, therefore, largely governed by the underlying classical transport. Remarkably, we will show below that the NESS still exhibits clear quantum correlations.

Quantum correlations—We consider the buildup of quantum correlations during the time evolution. The half-system negativity is defined as $\mathcal{N} = (||\rho^\Gamma||_1 - 1)/2$, where ρ^Γ denotes the partially transposed density matrix. $\mathcal{N}$ can witness entanglement via detection of violations of the positivity of the partial transpose ρ^Γ (PPT criterion) [72,73]. From Figs. 3(a)–3(c), we generally observe that the negativity initially increases sharply in all regions of the phase diagram. This is expected, as initially quantum correlations are generated through coherent transport. The behavior after the initial peak depends on the region of the phase diagram. In the MC phase, the negativity vanishes rapidly following this initial spike, and increasing system sizes do not support a longer-lived negativity. Conversely, in the LD and HD phase, as the system size increases, the initial negativity spike broadens, indicating that the corresponding entanglement survives for longer times when the system size is increased. However, in the finite-size settings, the negativity eventually dies out and does not play a role in the correlation structure of the

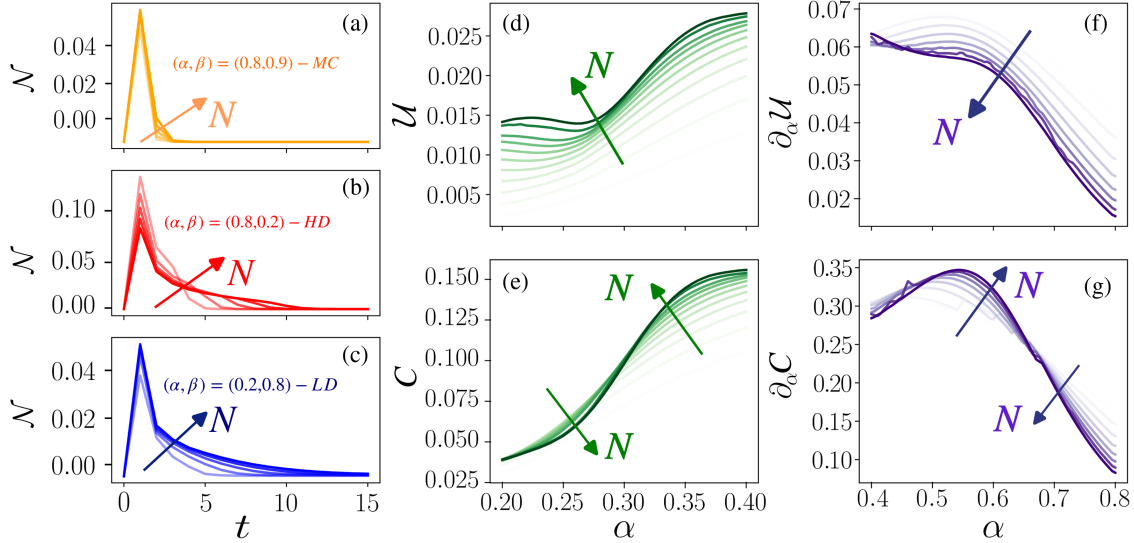


FIG. 3. Quantum correlations. (a)–(c) Time evolution of the half-system entanglement negativity for $N = 6, 8, 10, 12, 14$ in the MC, HD, and LD phase, respectively [see markers in Fig. 2(a)]. (d) and (e) show the maximum two-site LQU and coherence with respect to the central bulk qubit across the HD-LD transition ($\beta = 0.3$). (f) and (g) show the derivative of the same quantities across the LD-MC transitions ($\beta = 0.7$). All figures from (d)–(g) are plotted for $N = 6, 8, 10, 12, 14, 16, 18, 20, 22, 26, 30$.

finite-size NESS. We have also verified that all two-site reduced density matrix in the NESS are separable and that moment ratios of the partially transposed density matrix [74] of the NESS for larger system sizes do not violate the PPT criterion [56], further substantiating the claim that generally the finite-size NESS correlations are not dominated by bipartite entanglement. Our findings in the HD and LD phases indicate that a finite entanglement negativity may persist in the thermodynamic limit, because the time over which it survives increases with system size. This contrasts with prior observations of weak correlation signatures in strongly driven systems [28] and reveals that substantial quantum correlations may persist in the steady state.

Since we find that the negativity consistently vanishes in the long-time limit in our simulations, it raises the question whether other types of quantum correlations may be present in the NESS. Thus, we consider the local quantum uncertainty (LQU) $\mathcal{U}$ [75], which is closely related to quantum discord (see Ref. [76] for a review) and a coherence measure $\mathcal{C}$ in the classical basis that quantifies superpositions. To evaluate these quantities, we first pick a central bulk site at site $c = N/2$ and compute all two-site reduced density matrices $\rho_{j,c}$. We then evaluate $\mathcal{U}$ by performing local measurements on the central qubit and $\mathcal{C}$ by summing the absolute values of all off-diagonal elements of $\rho_{j,c}$ and consider their maximum across the phase boundaries. These measures are particularly appealing, because they are experimentally accessible via two-site quantum state tomography [77–80].

The results are shown in Figs. 3(d)–3(g). Across the LD-HD transition, we see that interestingly both $\mathcal{U}$ and $\mathcal{C}$

clearly witness the transition by an uptake in absolute value around the transition point; see Figs. 3(d) and 3(e). However, both phases do show extensive $\mathcal{U}$ for the system sizes we investigated, whereas $\mathcal{C}$ indicates a crossing in the vicinity of the coexistence line. Across the continuous LD-MC transition, $\partial_\alpha \mathcal{C}$ shows a crossing in the vicinity of the critical point, whereas the derivative $\partial_\alpha \mathcal{U}$ shows a strong drop upon entering the MC phase. We observed similar results across the remaining HD-MC transition.

These results indicate that local two-site superpositions in the bulk are most pronounced in the MC phase and more pronounced in the HD phase than the LD phase. This can be understood from the structure of the mixed state, whose dominant weight comes from pure states with nearest-neighbor empty-full pairings ($|\circ \bullet\rangle, |\bullet \circ\rangle$), implying that local clusters are typically in the single-particle sector. These states are mapped by the coherent transport U into local superpositions whereas the zero- and two-particle sectors are not mapped into superpositions. Therefore, the HD phase which has larger overlap with the single-particle sector than the LD phase also shows larger superpositions. This explains the steep signal of the coherence and the LQU across the LD-HD transition, since the typical two-site configurations change drastically across the coexistence line. Moreover, the MC phase is a homogeneous phase with the largest contribution from the single-particle sector and, thus, the most superpositions. This indicates that our quantum witnesses do flatten out in the MC phase, yielding a diminished derivative as can be seen in Fig. 3.

Thus, the LQU and the coherence qualitatively reproduce the phase diagram from Fig. 2. In Fig. 4, we show the results for $\mathcal{U}$ and $\mathcal{C}$ for $N = 30$ obtained from the matrix

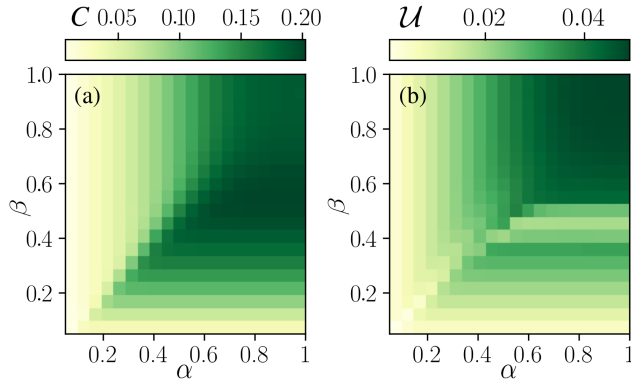


FIG. 4. Phase diagram from quantum correlations. (a) and (b) show the two-site the coherence and LQU, respectively, as described in the text, for the NESS that produced Fig. 2(a).

product states that also produced the occupation phase diagram in Fig. 2(a). We see that both witnesses reproduce the structure of the quantum phase diagram, indicating that quantum correlations beyond entanglement are an essential feature of the NESS and are capable of detecting non-equilibrium phase transitions.

Conclusion—We introduced a QCA model for coherent and stochastic transport and analyzed a quantum extension of the paradigmatic TASEP. While the quantum dynamics largely preserves the classical transport behavior at the level of local observables, leaving the phase diagram essentially unchanged, its correlation structure is fundamentally different. We established that entanglement plays a crucial role in the approach to the NESS, with its decay sharply distinguishing between the different phases. Strikingly, although no bipartite entanglement is detected in the NESS, the steady state exhibits clear quantum correlations beyond entanglement, including discord and coherence. It would be interesting to further explore how these correlations affect the broader thermodynamic properties of the NESS [81] and their impact on transport in realistic complex networks.

Our results show that nonclassical correlations persist in dissipative transport systems and even encode collective properties such as the phase structure in the absence of entanglement. This establishes dissipative quantum transport as a setting where all quantum correlations remain relevant. More broadly, our results suggest that strongly driven quantum systems generically host rich quantum correlation structures.

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Data availability—The data displayed in the figures are available [82].

- [1] C. T. MacDonald, J. H. Gibbs, and A. C. Pipkin, *Biopolymers* **6**, 1 (1968).
- [2] G. Lakatos and T. Chou, *J. Phys. A Math. Gen.* **36**, 2027 (2003).
- [3] T. Tripathi and D. Chowdhury, *Phys. Rev. E* **77**, 011921 (2008).
- [4] S. Scott and J. Szavits-Nossan, *Phys. Biol.* **17**, 015004 (2019).
- [5] J. Szavits-Nossan and R. Grima, *Phys. Rev. E* **108**, 034405 (2023).
- [6] I. Neri, N. Kern, and A. Parmeggiani, *Phys. Rev. Lett.* **107**, 068702 (2011).
- [7] N. Bunzarova, N. Pesheva, and J. Brankov, *Phys. Rev. E* **89**, 032125 (2014).
- [8] K. Nagel, *Phys. Rev. E* **53**, 4655 (1996).
- [9] P. Bonnin, I. Stansfield, M. C. Romano, and N. Kern, *Phys. Rev. E* **105**, 034117 (2022).
- [10] S. Prolhac, *Phys. Rev. Lett.* **116**, 090601 (2016).
- [11] D. A. Adams, R. K. P. Zia, and B. Schmittmann, *Phys. Rev. Lett.* **99**, 020601 (2007).
- [12] M. Gorissen, A. Lazarescu, K. Mallick, and C. Vanderzande, *Phys. Rev. Lett.* **109**, 170601 (2012).
- [13] J. de Gier and F. H. L. Essler, *Phys. Rev. Lett.* **107**, 010602 (2011).
- [14] M. Depken and R. Stinchcombe, *Phys. Rev. Lett.* **93**, 040602 (2004).
- [15] J. Brankov, N. Pesheva, and N. Valkov, *Phys. Rev. E* **61**, 2300 (2000).
- [16] N. Z. Bunzarova, N. C. Pesheva, and A. M. Povolotsky, *Phys. Rev. E* **109**, 044132 (2024).
- [17] R. A. Blythe and M. R. Evans, *J. Phys. A Math. Theor.* **40**, R333 (2007).
- [18] B. Derrida, M. R. Evans, V. Hakim, and V. Pasquier, *J. Phys. A Math. Gen.* **26**, 1493 (1993).
- [19] B. Derrida, *Phys. Rep.* **301**, 65 (1998).
- [20] G. M. Schütz, *J. Stat. Phys.* **88**, 427 (1997).
- [21] O. Golinelli and K. Mallick, *J. Phys. A Math. Gen.* **39**, 12679 (2006).
- [22] A. Rákos and G. Schütz, *J. Stat. Phys.* **118**, 511 (2005).
- [23] H. Spohn, *Physica (Amsterdam)* **369A**, 71 (2006).
- [24] N. Rajewsky and M. Schreckenberg, *Physica (Amsterdam)* **245A**, 139 (1997).
- [25] T. Chou, K. Mallick, and R. K. P. Zia, *Rep. Prog. Phys.* **74**, 116601 (2011).
- [26] D. Helbing, *Rev. Mod. Phys.* **73**, 1067 (2001).
- [27] K. Kavanagh, S. Dooley, J. K. Slingerland, and G. Kells, *New J. Phys.* **24**, 023024 (2022).
- [28] K. Temme, M. M. Wolf, and F. Verstraete, *New J. Phys.* **14**, 075004 (2012).
- [29] S. Wolfram, *Rev. Mod. Phys.* **55**, 601 (1983).
- [30] O. Martin, A. M. Odlyzko, and S. Wolfram, *Commun. Math. Phys.* **93**, 219 (1984).
- [31] M. Gardner, *Sci. Am.* **223**, 120 (1970).
- [32] S. Wolfram, *Phys. Rev. Lett.* **54**, 735 (1985).
- [33] T. Farrelly, *Quantum* **4**, 368 (2020).
- [34] R. White, V. Ramesh, A. Impertro, S. Anand, F. Cesa, G. Giudici, T. Iadecola, H. Pichler, and H. Bernien, *arXiv:2601.16257*.
- [35] F. Cesa, A. D. Fini, D. A. Korbany, R. Tricarico, H. Bernien, H. Pichler, and L. Piroli, *arXiv:2601.16961*.

- [36] O. Duranthon and G. Di Molfetta, *Phys. Rev. A* **103**, 032224 (2021).
- [37] E. Gillman, F. Carollo, and I. Lesanovsky, *Phys. Rev. Lett.* **125**, 100403 (2020).
- [38] K. Sfairopoulos, L. Causer, J. F. Mair, and J. P. Garrahan, *Phys. Rev. E* **113**, 064112 (2025).
- [39] I. Lesanovsky, K. Macieszczak, and J. P. Garrahan, *Quantum Sci. Technol.* **4**, 02LT02 (2019).
- [40] T. M. Wintermantel, Y. Wang, G. Lochead, S. Shevate, G. K. Brennen, and S. Whitlock, *Phys. Rev. Lett.* **124**, 070503 (2020).
- [41] E. Gillman, F. Carollo, and I. Lesanovsky, *Phys. Rev. E* **107**, L022102 (2023).
- [42] M. Boneberg, F. Carollo, and I. Lesanovsky, *Phys. Rev. A* **111**, 062405 (2025).
- [43] M. Boneberg, S. Kochsiek, G. Peretto, and I. Lesanovsky, *arXiv:2509.18262* [Phys. Rev. A (to be published)].
- [44] S. Wald, L. H. Yao, T. Platini, C. Hooley, and F. Carollo, *Quantum* **9**, 1742 (2025).
- [45] M. Albert, D. Bernard, T. Jin, S. Scopa, and S. Wei, *arXiv:2601.16883*.
- [46] D. Bernard, *J. Phys. A Math. Theor.* **54**, 433001 (2021).
- [47] D. Bernard, F. H. L. Essler, L. Hruza, and M. Medenjak, *SciPost Phys.* **12**, 042 (2022).
- [48] D. Bernard, T. Jin, S. Scopa, and S. Wei, *Phys. Rev. E* **112**, 034106 (2025).
- [49] V. Alba, *Phys. Rev. B* **113**, 075145.
- [50] A. Russotto, F. Ares, P. Calabrese, and V. Alba, *J. Stat. Mech.* (2026) 033103.
- [51] R. Khasseh, S. Wald, R. Moessner, C. A. Weber, and M. Heyl, *Phys. Rev. Lett.* **135**, 248302 (2025).
- [52] H. Labuhn, D. Barredo, S. Ravets, S. de Léséleuc, T. Macrì, T. Lahaye, and A. Browaeys, *Nature (London)* **534**, 667 (2016).
- [53] H. Bernien, S. Schwartz, A. Keesling, H. Levine, A. Omran, H. Pichler, S. Choi, A. S. Zibrov, M. Endres, M. Greiner, V. Vuletić, and M. D. Lukin, *Nature (London)* **551**, 579 (2017).
- [54] A. Browaeys and T. Lahaye, *Nat. Phys.* **16**, 132 (2020).
- [55] S. Ebadi, T. T. Wang, H. Levine, A. Keesling, G. Semeghini, A. Omran, D. Bluvstein, R. Samajdar, H. Pichler, W. W. Ho, S. Choi, S. Sachdev, M. Greiner, V. Vuletić, and M. D. Lukin, *Nature (London)* **595**, 227 (2021).
- [56] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/2cvx-6hqz> for the derivation of the Kraus map, the solution of the classical TASEP, the continuous-time limit, additional analyses of the nonequilibrium phase diagram, and supplementary numerical results, which includes Refs. [57–59].
- [57] T. Prosen, *Phys. Rev. Lett.* **107**, 137201 (2011).
- [58] A. Elben, R. Kueng, H.-Y. R. Huang, R. van Bijnen, C. Kokail, M. Dalmonte, P. Calabrese, B. Kraus, J. Preskill, P. Zoller, and B. Vermersch, *Phys. Rev. Lett.* **125**, 200501 (2020).
- [59] R. Horodecki, P. Horodecki, M. Horodecki, and K. Horodecki, *Rev. Mod. Phys.* **81**, 865 (2009).
- [60] D. Karevski and T. Platini, *Phys. Rev. Lett.* **102**, 207207 (2009).
- [61] E. Gillman, F. Carollo, and I. Lesanovsky, *Phys. Rev. Lett.* **127**, 230502 (2021).
- [62] E. Gillman, F. Carollo, and I. Lesanovsky, *Phys. Rev. E* **106**, L032103 (2022).
- [63] M. Boneberg, F. Carollo, and I. Lesanovsky, *New J. Phys.* **25**, 093020 (2023).
- [64] As U does not act on the ancilla, we discard its label here.
- [65] D. Karevski, V. Popkov, and G. M. Schütz, *Phys. Rev. Lett.* **110**, 047201 (2013).
- [66] F. Carollo, J. P. Garrahan, and I. Lesanovsky, *Phys. Rev. B* **98**, 094301 (2018).
- [67] M. Žnidarič, *Phys. Rev. Lett.* **112**, 040602 (2014).
- [68] M. Žnidarič, *Phys. Rev. B* **90**, 115156 (2014).
- [69] E. Ilievski and B. Žunkovič, *J. Stat. Mech.* (2014) P01001.
- [70] G. Vidal, *Phys. Rev. Lett.* **93**, 040502 (2004).
- [71] G. Vidal, *Phys. Rev. Lett.* **98**, 070201 (2007).
- [72] G. Vidal and R. F. Werner, *Phys. Rev. A* **65**, 032314 (2002).
- [73] M. B. Plenio, *Phys. Rev. Lett.* **95**, 090503 (2005).
- [74] J. Carrasco, M. Votto, V. Vitale, C. Kokail, A. Neven, P. Zoller, B. Vermersch, and B. Kraus, *Phys. Rev. A* **109**, 012422 (2024).
- [75] D. Girolami, T. Tufarelli, and G. Adesso, *Phys. Rev. Lett.* **110**, 240402 (2013).
- [76] A. Bera, T. Das, D. Sadhukhan, S. Singha Roy, A. Sen(De), and U. Sen, *Rep. Prog. Phys.* **81**, 024001 (2017).
- [77] M. Paris and J. Rehacek, *Quantum State Estimation* (Springer, Berlin Heidelberg, 2004).
- [78] R. Auccaise, J. Maziero, L. C. Céleri, D. O. Soares-Pinto, E. R. deAzevedo, T. J. Bonagamba, R. S. Sarthour, I. S. Oliveira, and R. M. Serra, *Phys. Rev. Lett.* **107**, 070501 (2011).
- [79] K. Kim and J. Ahn, *PRX Quantum* **4**, 020316 (2023).
- [80] M. Rambach, M. Qaryan, M. Kewming, C. Ferrie, A. G. White, and J. Romero, *Phys. Rev. Lett.* **126**, 100402 (2021).
- [81] P. Strasberg, G. Schaller, T. Brandes, and M. Esposito, *Phys. Rev. X* **7**, 021003 (2017).
- [82] U. Sen, F. Carollo, and S. Wald, [10.5281/zenodo.21377919](https://arxiv.org/abs/10.5281/zenodo.21377919) (2026).