

Active Quantum Matter from Monitored Pure-State Dynamics

Jacob F. Steiner¹, Felix von Oppen², and Reinhold Egger³

¹*Department of Physics and Institute for Quantum Information and Matter, California Institute of Technology, Pasadena, California 91125, USA*

²*Dahlem Center for Complex Quantum Systems, Fachbereich Physik, and Halle-Berlin-Regensburg Cluster of Excellence CCE, Freie Universität Berlin, 14195 Berlin, Germany*

³*Institut für Theoretische Physik, Heinrich-Heine-Universität, D-40225 Düsseldorf, Germany*



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Quantum many-body systems coupled to out-of-equilibrium reservoirs can behave as active matter and exhibit signs of flocking. However, the resulting steady states are highly mixed and carry only weak quantum signatures. We show that signatures of active matter also arise in ensembles of pure states undergoing monitored quantum dynamics. We consider a spinful Luttinger liquid subject to measurement processes that shuffle spin-up particles to the left and spin-down particles to the right. For weak monitoring strengths and ferromagnetic spin interactions, we find power-law quantum correlations between spin density and charge current, which we identify as a hallmark of active quantum matter. The monitoring plays a dual role, generating the quantum active correlations for weak strengths while driving a Berezinskii-Kosterlitz-Thouless phase transition to a short-range correlated state at larger strengths.

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Introduction—Classical active matter has attracted enormous attention due to its rich phenomenology and broad range of applicability [1–4]. A representative subclass of active matter are interacting assemblies of self-propelled particles which spontaneously align their direction of motion and form a flock [5–10]. The self-propulsion allows these systems to bypass the Mermin-Wagner theorem and exhibit spontaneous breaking of a continuous symmetry in two spatial dimensions [5,11]. Even one-dimensional (1D) active Ising chains, while not technically long-range ordered, exhibit a transition from a disordered to a phase-separated regime that is characterized by finite-life-time clusters of particles moving together on a disordered background [6,7,12].

It is natural to ask whether and how these phenomena translate to quantum systems. Very recently, first forays into this largely uncharted territory have been undertaken, broadly following two distinct approaches. First, a number of works aimed at establishing minimal structures required to endow a quantum particle with the ability to propel itself. These considered quantum walks [13], classical noise [14–17], or heat-to-motion conversion [18]. Active motion has also been investigated for defects in spin textures [19,20]. Second, collective effects of interacting

self-propelled quantum particles were studied based on nonhermitian Hamiltonian [14,21] and Lindbladian [22,23] dynamics. The models studied in Refs. [14,21,22] are closely related to the active Ising model of Ref. [6], describing spin- $\frac{1}{2}$ particles subject to incoherent directed hopping and spin alignment mechanisms. In particular, Khasseh *et al.* [22] predict signatures of flocking to survive in the presence of sufficiently weak quantum spin-flip processes.

These first steps toward a theory of active quantum matter leave key issues open. If directed motion and alignment are due to incoherent (and hence essentially classical) processes, one expects low-purity mixed states at long times, so quantum signatures are observable only as transient correlations. Here we pose, and answer in the affirmative, the following question: Can active quantum matter be based solely on the fundamental ingredients of quantum many-body dynamics, namely unitary time evolution and quantum measurements [24,25]? The active nature of the system then originates from quantum measurements. Importantly, this makes quantumness an essential ingredient as (weak) classical measurements do not influence the system dynamics and hence cannot induce directionality or self-propulsion. Thus, quantum measurements provide a general construction principle for active matter in the deep quantum limit. In fact, for a given measurement trajectory, the system remains in a pure quantum state at all times given a pure initial state. We show that quantum active behavior can be realized based on this principle. In addition to the rich and interesting physics revealed by our case study, this opens the door to exploring

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quantum mechanisms of activity in a wide range of systems of different dimensionalities, geometries, or quantum statistics, to mention just a few.

We study a 1D quantum active model of spin- $\frac{1}{2}$ fermions with ferromagnetic Ising coupling, similar to Ref. [22], but employ only monitored quantum dynamics, i.e., unitary dynamics combined with weak continuous measurements [24]. We design the measurements to generate directed hopping and thus self-propulsion. Quantum measurements allow energy and information exchange between individual particles and the detectors, thereby equipping our model with the defining characteristic of quantum active matter [10]. The resulting pure-state dynamics exhibit features of quantum active motion manifested in quasi-long-range (algebraically decaying) quantum correlations between spin density and charge current. Initially, these correlations grow with increasing monitoring strength or ferromagnetic coupling. In this sense, the system forms an incipient quantum flock.

Importantly, the monitoring plays a dual role. Beyond a critical strength, it drives a transition to a phase in which the correlations between spin density and charge current become short-ranged. This Berezinskii-Kosterlitz-Thouless (BKT) transition [26] between phases with distinct levels of quantum active correlations is directly related to the charge sharpening transition in U(1)-symmetric monitored quantum many-body systems [27–30]. We discuss our monitoring scheme—implemented via nonhermitian jump operators that drive self-propulsion—and its competition with the unitary dynamics within an intuitive lattice model. We use a corresponding bosonized continuum theory, a nonhermitian spinful Luttinger liquid with a sine-Gordon nonlinearity, to gain analytical access to the BKT transition. This builds on recent work on spinless models with monitoring of hermitian operators [31–33].

Lattice models of generic U(1)-symmetric monitored systems [27–30] exhibit a transition between volume- and area-law entangled phases in addition to the charge-sharpening transition [34–44]. Calculations indicate that this entanglement transition occurs at even stronger monitoring within the charge-sharp phase [27–30]. By analogy, we expect the BKT transition to separate two distinct volume-law entangled phases, while the entanglement transition might be viewed as the boundary between quantum and classical active phases.

Lattice model—We describe the monitored evolution of a 1D chain of spin- $\frac{1}{2}$ fermions $c_{x,\sigma}$ with site index $x \in \mathbb{Z}$ and spin $\sigma \in \{\uparrow, \downarrow\} \equiv \{+1, -1\}$ by the stochastic Schrödinger equation [24,25,45]:

$$i \frac{d}{dt} |\psi_c\rangle = \left\{ H - \frac{i}{2} \sum_{x,\sigma} [L_{x,\sigma}^\dagger L_{x,\sigma} - 2W_{x,\sigma} L_{x,\sigma} + W_{x,\sigma}^2 + 2\xi_{x,\sigma}(L_{x,\sigma} - W_{x,\sigma})] \right\} |\psi_c\rangle. \quad (1)$$

The subscript c denotes conditioning on the measurement trajectory, and the terms involving $W_{x,\sigma} = \text{Re}\langle\psi_c|L_{x,\sigma}|\psi_c\rangle$ ensure normalization of $|\psi_c\rangle$. The unitary dynamics obey a Hamiltonian H including nearest-neighbor hopping t_0 and ferromagnetic Ising exchange $J_z > 0$,

$$H = \sum_x [-t_0(c_x^\dagger c_{x+1} + \text{H.c.}) - J_z S_x^z S_{x+1}^z] + H_{\text{flip}}, \quad (2)$$

where $S_x^z = \frac{1}{2} c_x^\dagger \sigma_3 c_x$ and $c_x = [c_{x,\uparrow}, c_{x,\downarrow}]^\top$. The Pauli matrices ($\sigma_1, \sigma_2, \sigma_3$) act in spin space. We also allow for spin-nonconserving dynamics through a time-reversal odd Zeeman term, $H_{\text{flip}} = h \sum_x c_x^\dagger \sigma_1 c_x$, or a time-reversal even Rashba spin-orbit coupling, $H_{\text{flip}} = -i\alpha \sum_x c_x^\dagger \sigma_1 c_{x+1} + \text{H.c.}$ The self-propulsion of spin- σ fermions is achieved by monitoring with the nonhermitian jump operators,

$$L_{x,\sigma} = \sqrt{\gamma} c_{x+\sigma,\sigma}^\dagger c_{x,\sigma}, \quad (3)$$

with monitoring strength γ . Physically, we continuously monitor jumps of a spin- σ particle by one site in a spin-dependent direction, $x \rightarrow x + \sigma$. The random sequence of measurement outcomes is encoded by the Gaussian white noise sources $\xi_{x,\sigma}$ with $\overline{\xi_{x,\sigma}(t)} = 0$ and correlations $\overline{\xi_{x,\sigma}(t) \xi_{x',\sigma'}(t')} = \delta_{xx'} \delta_{\sigma\sigma'} \delta(t - t')$, where the overline denotes averaging over the ensemble of measurement outcomes. We emphasize that, for a single measurement trajectory, Eq. (1) preserves purity at all times t .

The kinetic processes of our model are graphically represented in Fig. 1. Spin-independent coherent hopping without preferred direction (amplitude t_0) builds up entanglement. Spin-dependent directional stochastic jumps are induced by monitoring [jump operators in Eq. (3)]. The monitoring detects jumps of spin-dependent directionality. In principle, this can be realized experimentally, e.g., by exploiting cotunneling by coupling the sites to an auxiliary chain of spin-polarized detector quantum dots [46] and measuring their charge state, as explained in detail in Supplemental Material (SM) [47]. Our model differs from the one discussed in Ref. [22] even with regard to the Lindblad evolution $\dot{\rho}_1 = \mathcal{L}\rho_1 = -i[H, \rho_1] + \sum_{x,\sigma} (L\rho_1 L^\dagger - \frac{1}{2}\{L^\dagger L, \rho_1\})$ of the ensemble-averaged state

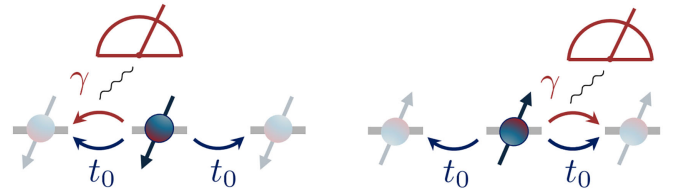


FIG. 1. Cartoon of kinetic processes in the 1D monitored lattice model defined by Eqs. (1)–(3). Spin-independent and bidirectional motion is due to hopping with amplitude t_0 (blue). Spin-dependent unidirectional motion is induced by monitoring with monitoring strength γ (red).

$\rho_1 = |\psi_c\rangle\langle\psi_c|$ [24,25,48]. First, we include spin-independent coherent hopping. Second, the Ising coupling enters the unitary dynamics. Third, we do not include an explicit alignment process. As a consequence, ρ_1 heats up to become completely mixed, $\rho_1(t \rightarrow \infty) \propto \mathbb{1}$, suppressing signatures of quantum active motion in the ensemble-averaged steady state [49].

Importantly, signatures of quantum active motion persist at long times in the pure-state dynamics generated by Eq. (1). We consider measurement-averaged equal-time connected correlation functions,

$$C_{AB}(x) = \overline{\langle\psi_c|A_x B_0|\psi_c\rangle} - \overline{\langle\psi_c|A_x|\psi_c\rangle} \overline{\langle\psi_c|B_0|\psi_c\rangle}. \quad (4)$$

These exhibit nontrivial behavior due to the second term, which is nonlinear in the system's density matrix $|\psi_c\rangle\langle\psi_c|$ [32]. Here, A_x and B_0 represent local operators (charge and spin densities as well as currents) at sites x and 0 , respectively. We express $C_{AB}(x)$ in terms of the two-replica density matrix averaged over measurement outcomes, $\rho_2 = |\psi_c\rangle\langle\psi_c| \otimes |\psi_c\rangle\langle\psi_c|$. This yields

$$C_{AB}(x) = \frac{1}{2} \left\langle (A_x^{(1)} - A_x^{(2)})(B_0^{(1)} - B_0^{(2)}) \right\rangle, \quad (5)$$

where the superscript indicates action on replica $i = 1, 2$ and $\langle \dots \rangle = \text{tr}(\dots \rho_2)$.

A natural signature for active self-propulsion is offered by the correlations between the local spin density (“magnetization”) q_s and the charge current density j_c . In particular, such correlations are significant if spin-up

(spin-down) particles move to the right (left) since a spin-density imbalance is then accompanied by a charge current density. We find that, below a critical measurement rate, the corresponding steady-state correlations exhibit power-law behavior at long distances, $C_{\rho_s j_c}(x) \sim -\Delta_{sc}/x^2$. We emphasize that these are quantum correlations of a pure-state ensemble and are thus accompanied by significant entanglement. The coefficient Δ_{sc} is an increasing function of γ : activity is driven by the measurements. When γ exceeds a critical strength, the pure-state dynamics undergoes a transition into a phase with exponentially decaying $C_{\rho_s j_c}(x)$. This is summarized in the phase diagram in Fig. 2(a), where g_s is the Luttinger liquid parameter quantifying the Ising interaction (see below). The color scale indicates the active correlations $\propto \Delta_{sc}$, which increase with γ right up to the transition; see Fig. 2(b) for a linecut.

Bosonization—Rather than working directly with the lattice model in Eqs. (2) and (3), we consider a bosonized continuum theory, $H \rightarrow H_0 + \delta H$ and $L_{x,\sigma} \rightarrow L_\sigma(x)$, which captures this BKT transition and allows for analytical progress. (From now on, x is a continuous coordinate.) For generic incommensurate band fillings, the quadratic Hamiltonian,

$$H_0 = \frac{v_F}{2} \sum_\nu \int dx \left[(\partial_x \theta_\nu)^2 + \frac{1}{g_\nu^2} (\partial_x \varphi_\nu)^2 \right], \quad (6)$$

describes the dynamics of the charge ($\nu = c$) and spin ($\nu = s$) degrees of freedom in terms of two Luttinger liquids with Fermi velocity v_F . Here, $\partial_x \theta_\nu(x)$ and $\varphi_\nu(x)$ are canonically conjugate real bosonic fields

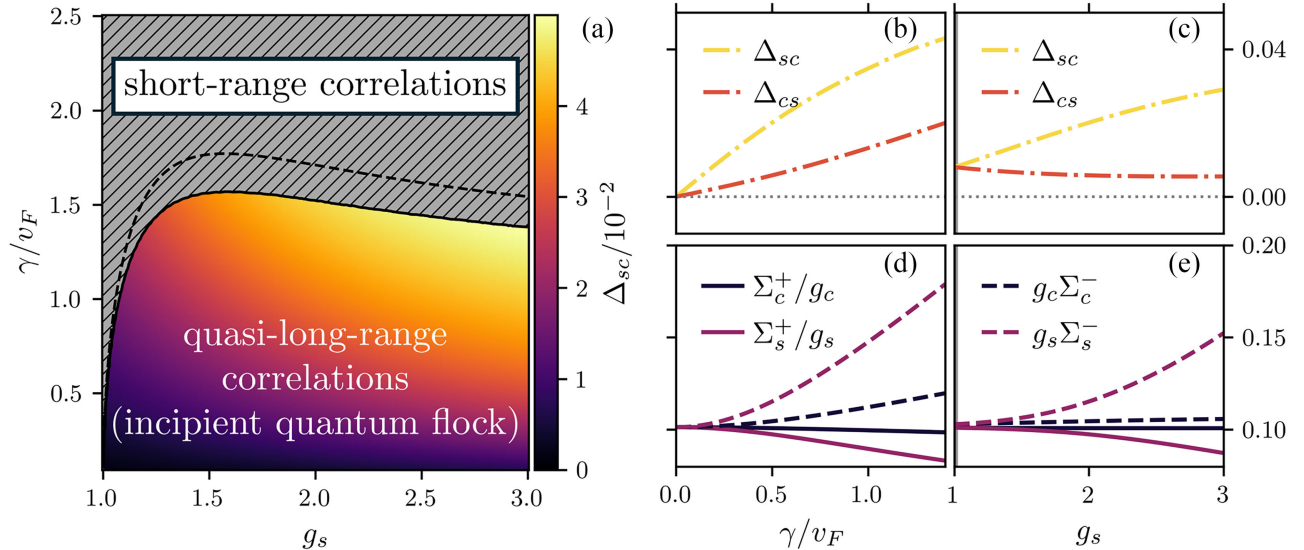


FIG. 2. (a) Phase diagram in the g_s - γ plane obtained from the RG equations (13) with $g_c = 1$ and $\tilde{c} = \tilde{s} = 1/\sqrt{2}$; see Eq. (12). The quasi-long-range quantum active phase (small γ) and the short-range phase (large γ) are separated by a BKT transition (black line). In the quantum active phase, color indicates the magnitude of Δ_{cs} . The dashed line is an analytical estimate for the phase boundary (see main text). (b)–(e) Dependence of the coefficients $\{\Sigma_{c/s}^\pm, \Delta_{cs/sc}\}$ on γ and g_s , which determine the correlations at the fixed point $\lambda = 0$, see Eq. (15), using $g_s = 2$ in (b) and (d) and $\gamma/v_F = 0.5$ in (c) and (e).

with $[\varphi_\nu(x), \partial_{x'}\theta_\nu(x')] = i\delta_{\nu\nu'}\delta(x-x')$ [26,50,51]. Their physical meaning is explained after Eq. (7); see also Ref. [26]. Measuring lengths in units of the lattice spacing, the spin sector has the Luttinger parameter $g_s \simeq 1/[1 - J_z/2\pi v_F]^{1/2} > 1$, while $g_c = 1$ in the charge sector. Zeeman and Rashba terms as well as backscattering processes due to J_z give rise to sine-Gordon terms collected in δH (see SM [47]). Similarly, the bosonized measurement operators $L_\sigma(x) = M_\sigma(x) + iK_\sigma(x)$ with hermitian M_σ and K_σ decompose into linear and sine-Gordon terms, $M_\sigma = M_{0,\sigma} + \delta M_\sigma$ and $K_\sigma = K_{0,\sigma} + \delta K_\sigma$. Here,

$$M_{0,\sigma} = \sqrt{\frac{\gamma}{2\pi}} \tilde{c} \partial_x(\varphi_c + \sigma\varphi_s) \sim \text{density}_\sigma, \quad (7a)$$

$$K_{0,\sigma} = \sqrt{\frac{\gamma}{2\pi}} \tilde{s} \sigma \partial_x(\theta_c + \sigma\theta_s) \sim \sigma \times \text{current}_\sigma \quad (7b)$$

stem from momentum-conserving processes. ($\tilde{c}$ and $\tilde{s}$ with $\tilde{c}^2 + \tilde{s}^2 = 1$ are real constants controlled by the band filling.) The hermitian part $M_{0,\sigma}$ depends on the smooth component $\partial_x(\varphi_c + \sigma\varphi_s)/\sqrt{2}$ of the density of spin- σ particles. Similarly, the antihermitian part $iK_{0,\sigma}$ depends on the smooth component $\sigma\partial_x(\theta_c + \sigma\theta_s)/\sqrt{2}$ of the Hamiltonian current density in the σ direction due to spin- σ particles. Further, $2k_F$ processes give rise to the sine-Gordon contributions,

$$\delta M_\sigma = \sqrt{\frac{\gamma}{\pi^2}} \sigma \sin[2k_F x + \sqrt{2\pi}(\varphi_c + \sigma\varphi_s)], \quad (8)$$

while $\delta K_\sigma = 0$. Because of the large difference in momentum transfers, we can treat the $2k_F$ term as a separate measurement process [32]. Thus, there are two sets of jump operators, $L_{0,\sigma} = M_{0,\sigma} + iK_{0,\sigma}$ and $\delta L_\sigma = \delta M_\sigma$. In the bosonized theory, heating of ρ_1 manifests as unbounded growth of the correlation functions $\text{tr}[\chi_{\pm,\nu}(x)\chi_{\pm,\nu'}(x)\rho_1(t)]$ for $t \rightarrow \infty$, where $\chi_{+,\nu} = \varphi_\nu$ and $\chi_{-,\nu} = \theta_\nu$ [31]. We recall that ρ_1 is the ensemble-averaged state while ρ_2 refers to the averaged two-replica density matrix.

Nonhermitian sine-Gordon theory—We now consider the time evolution of ρ_2 . Correlations of charge and spin densities and/or currents can be traced to $C_{AB}(x)$ with $A, B \sim \partial_x\chi_{\pm,\nu}$, which depends only on the relative modes $\chi^{(r)} = (\chi^{(1)} - \chi^{(2)})/\sqrt{2}$; see Eq. (5). Conversely, as a consequence of the heating of ρ_1 , correlations of the center-of-mass (absolute) modes $\chi^{(a)} = (\chi^{(1)} + \chi^{(2)})/\sqrt{2}$ grow without bound for $t \rightarrow \infty$. In fact, symmetry under replica exchange gives $\langle \chi^{(a)} \rangle = \sqrt{2}\text{tr}(\chi\rho_1)$, $\langle \chi^{(r)} \rangle = 0$, and

$$\langle \chi_\alpha^{(a)} \chi_\beta^{(a)} \rangle = 2\text{tr}(\chi_\alpha \chi_\beta \rho_1) - \langle \chi_\alpha^{(r)} \chi_\beta^{(r)} \rangle. \quad (9)$$

Assuming that correlations of the relative fields remain finite (as confirmed *a posteriori*), unbounded heating of ρ_1 translates to unbounded heating of correlations in the

absolute sector: $\langle \chi_\alpha^{(a)}(x)\chi_\beta^{(a)}(x) \rangle \rightarrow \infty$ and $\langle \chi^{(a)}(x) \rangle = 0$. To arrive at an effective description of the relative sector, we trace out the absolute modes; see SM [47] for details. Importantly, at long times, all contributions from δH become irrelevant under this procedure. Using $\rho_2^{(r)} = \text{tr}^{(a)}\rho_2$, the dynamics in the relative replica sector is then governed by

$$\dot{\rho}_2^{(r)} = -iH_{\text{eff}}^{(r)}\rho_2^{(r)} + i\rho_2^{(r)}[H_{\text{eff}}^{(r)}]^\dagger. \quad (10)$$

Mostly dropping the superscript (r) from now on, the effective nonhermitian sine-Gordon Hamiltonian $H_{\text{eff}} = v_F \int dx(\mathcal{H}_{\text{eff},0} + \delta\mathcal{H}_{\text{eff}})$ is given by

$$\mathcal{H}_{\text{eff},0} = \frac{1}{2} \sum_{\nu=c,s} [(\partial_x\theta_\nu)^2 + (k_\nu + \kappa^2)(\partial_x\varphi_\nu)^2] + \kappa(\partial_x\varphi_c\partial_x\theta_s + \partial_x\varphi_s\partial_x\theta_c), \quad (11a)$$

$$\delta\mathcal{H}_{\text{eff}} = i\lambda \cos(\sqrt{4\pi}\varphi_c) \cos(\sqrt{4\pi}\varphi_s), \quad (11b)$$

where we introduced the couplings

$$k_\nu = \frac{1}{g_\nu^2} - \frac{2i\tilde{c}^2\gamma}{\pi v_F} - \kappa^2, \quad \kappa = \frac{\tilde{c}\tilde{s}\gamma}{\pi v_F}, \quad \lambda = \frac{\gamma}{\pi^2 v_F}. \quad (12)$$

Note that γ enters both the parity-odd spin-charge coupling κ (giving rise to $\Delta_{sc} \neq 0$) and the parity-even nonlinearity λ (driving the BKT transition), highlighting the dual role of monitoring for quantum active matter. Importantly, Eq. (10) does not feature a jump term and thus describes the nonhermitian evolution of a pure state, $\rho_2^{(r)} = |\psi\rangle\langle\psi|$, toward a dark state $|\psi(t \rightarrow \infty)\rangle = |0\rangle$. Steady-state connected correlation functions are then given by $C_{AB}(x) = \frac{1}{2}\langle 0|A^{(r)}(x)B^{(r)}(0)|0\rangle$.

Renormalization and phase diagram—To find $|0\rangle$, we treat the nonlinear term $\propto \lambda$ in Eq. (11) using a perturbative renormalization group (RG) analysis [26,32]. The one-loop RG equations for the complex couplings (12) are (for a derivation, see SM [47])

$$\frac{d\lambda}{d\ell} = \left(2 - \sum_\nu \frac{1}{\sqrt{k_\nu}}\right)\lambda, \quad \frac{dk_\nu}{d\ell} = -f\lambda^2, \quad (13)$$

where ℓ is the standard RG flow parameter, $|f| \sim \mathcal{O}(1)$ depends weakly on k_ν , and κ is not renormalized. Solving Eq. (13) with the initial values (12) gives the phase diagram in Fig. 2(a), with the solid curve indicating the transition between a phase with quasi-long-range quantum active correlations, where $\lambda(\ell \rightarrow \infty) = 0$, and a phase with short-range correlations, where λ flows toward strong coupling. The RG equations (13) are characteristic for a BKT transition [26,32,33]. Compared to the 2D classical case, the BKT transition occurs here in a $(1+1)$ -dimensional

quantum system. In the quantum active phase, one may assume that k_ν remains unrenormalized since λ quickly approaches zero. The resulting estimate for the phase boundary, $\sum_\nu 1/\sqrt{k_\nu} = 2$, yields the dashed curve in Fig. 2(a). We find that the quantum active phase disappears without interactions ($g_s \rightarrow 1$).

Gaussian fixed points—We now study the steady-state correlations $C_{AB}(x)$ for all combinations of $A, B \in \{q_{c/s}, j_{c/s}\}$ at the $|\lambda| \rightarrow 0, \infty$ fixed points. Since these limits correspond to Gaussian fixed points, we obtain exact results for the long-distance behavior. With the mean density q_0 , charge and spin densities are given by $q_\nu = \sqrt{2/\pi} \partial_x \varphi_\nu + q_0 \delta_{\nu,c}$. The respective current densities j_ν include measurement-induced contributions $\propto \kappa$, see Refs. [47,52–54],

$$j_\nu = -\sqrt{2/\pi} v_F [\partial_x \theta_\nu - \kappa \partial_x \varphi_\nu] + v_F \kappa q_0 \delta_{\nu,s}, \quad (14)$$

with $\bar{\nu} = s, c$ for $\nu = c, s$. Note that monitoring induces a constant spin-dependent current. Consider first the case $\lambda \rightarrow 0$, where we find

$$\begin{pmatrix} C_{q_\nu q_{\bar{\nu}}} & C_{q_\nu j_{\bar{\nu}}} \\ C_{j_\nu q_{\bar{\nu}}} & C_{j_\nu j_{\bar{\nu}}} \end{pmatrix} \simeq -\frac{1}{x^2} \begin{pmatrix} \delta_{\nu\bar{\nu}} \Sigma_\nu^+ & v_F \delta_{\nu\bar{\nu}} \Delta_{\nu\bar{\nu}} \\ v_F \delta_{\nu\bar{\nu}} \Delta_{\bar{\nu}\nu} & v_F^2 \delta_{\nu\bar{\nu}} \Sigma_\nu^- \end{pmatrix}. \quad (15)$$

The six real coefficients $\{\Sigma_{c/s}^\pm, \Delta_{cs/sc}\}$ are illustrated as functions of γ and/or g_s in Figs. 2(b)–2(e). As in equilibrium Luttinger liquids [26], density-density and current-current correlations decay $\propto x^{-2}$, implying that members of the pure-state ensemble are significantly entangled [32,55]. Note that spin current (density) correlations are enhanced (reduced) by finite γ and by $g_s > 1$, while charge correlations are not strongly affected. Equation (15) reveals quantum active correlations between spin (charge) densities and charge (spin) currents from the $\Delta_{sc(cs)}$ terms. Notably, Δ_{sc} is enhanced by the Ising interaction while Δ_{cs} is weakened, which can be rationalized by noting that Ising interactions favor alignment of adjacent spins and thus quantum flocking. Finally, consider the strong-coupling phase with large $|\lambda|$. Here, the nonlinearity in Eq. (11) leads to a Gaussian fixed point with a mass term in the boson dispersion (mass m_λ). We then find exponential decay, $C_{A_\nu B_{\bar{\nu}}}(x) \propto e^{-|x|/\xi_{\nu\bar{\nu}}}$, with the length scale [47]

$$\xi_{\nu\bar{\nu}} \sim \frac{v_F}{|m_\lambda|} \max(\sqrt{|k_\nu|}, \sqrt{|k_{\bar{\nu}}|}). \quad (16)$$

We thus encounter only short-range correlations.

Discussion—We showed that quantum active matter can be realized in monitored interacting fermion chains. Monitoring causes self-propulsion but eventually also drives a BKT transition between phases with algebraic and short-ranged quantum active correlations. The algebraic phase is strongly quantum in that typical pure states

cannot be prepared by shallow circuits. The short-range correlated phase arises from the quantum Zeno effect under frequent measurements. This phase is expected to host a second transition to a weakly quantum area-law entangled phase [27–30], with typical pure states preparable by shallow circuits [56]. An explicit study of the entanglement transition is left for future work as how to describe it within the bosonization approach is currently unknown [33]. Matrix product state methods may be useful in this context. Experimental studies of the quantum active correlations require many copies of the ensemble of measurement trajectories. To deal with the associated postselection problem [42], one can use methods such as classical-quantum cross-correlations [57], classical postprocessing [58], decoding [59], space-time duality [60], or tree-like measurements [61]. The entanglement transition was successfully observed in this way [62–64]. We expect our work to stimulate theoretical and experimental research on active matter in the deep quantum limit by employing quantum measurements.

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Data availability—The raw data underlying the figures presented in this work are available at Zenodo [65]. Reference [65] contains the Zenodo URL site.

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