

Absorbing state phase transitions beyond directed percolation in dissipative quantum state preparation

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We show that absorbing state phase transitions where the absorbing state itself exhibits long-range phase coherence can lead to critical behavior distinct from directed percolation. To do this, we investigate a simple, purely dissipative quantum reaction-diffusion model, which may also be viewed as a dissipative quantum state preparation procedure for the (generalized) W state with errors. The “error” Lindblad jump operators preserve the W state as a dark state, but nonetheless act to decohere the system and induce the phase transition. We find cases where the preparation protocol is either fragile or robust against weak error quantum jump rates and show that local remnants of the coherence persist in the decohering phase. The distinct critical behavior stems from the spreading of coherence throughout the system at the critical point.

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I. INTRODUCTION

The dynamics of nonequilibrium quantum many-body systems exhibit striking phenomena which would be rare, if not impossible, in classical or equilibrium settings. Interest in this area has been especially spurred by recent experimental advances—across a variety of platforms—in the control of quantum systems at the atom-by-atom (or qubit-by-qubit) level [1–5].

Often, evolution in these systems has been viewed as a competition between the coherent Hamiltonian dynamics which may contribute to the quantum nature of a system, and the decohering effects of coupling to the environment or measurements which destroy entanglement and other quantum signatures. However, even in systems with no Hamiltonian dynamics, it is possible to generate entanglement solely through dissipation [6,7]. Additionally, the irreversible nature of dissipation allows for the one-way cooling of quantum systems to a desired steady state as well as provides the possibility for interactive control of a system through feedback. These properties—which allow dissipative dynamics to be especially robust to perturbation—make it a key tool in quantum information [8], showing up in quantum state preparation and stabilization [9–17], error correction [18–22], and as distinct models of universal quantum computation [6,23].

In condensed matter, the same experimental advancements in engineered coupling to the environment and quantum control have led to the discovery of new phases of matter and dynamical phase transitions which are unique to the nonequilibrium setting [24–29]. Nonequilibrium phase transitions are present in a wide variety of contexts [30–35]. An important class of systems which exhibit such transitions are

reaction-diffusion models, where particles (or reactants) move through a system by diffusion and react when they meet. Here, the system may exhibit phase transitions to absorbing states (i.e., configurations which may be approached by the dynamics, but not left); such transitions typically belong to the directed percolation (DP) universality class [36]. There has been a flurry of recent interest—on both theoretical [37–48] and experimental [49–52] fronts—into quantum extensions to reaction-diffusion models, with continuing debate and many open questions surrounding precisely when and how quantum effects may induce novel critical behavior distinct from DP.

In this paper, we show that critical behavior beyond DP can occur when the absorbing state itself has long range coherence. This distinct behavior stems from dynamics associated with the spreading of quantum coherence throughout the system at the critical point. To show these results, we provide a prototypical example of a quantum reaction-diffusion model with a transition to a long range coherent absorbing state. We find, furthermore, this coherence may be induced entirely through dissipative means, with no Hamiltonian evolution. The model may also be viewed as a dissipative quantum state preparation procedure—preparing the long range entangled absorbing state—with additional “error” environmental couplings that ruin the preparation protocol. The absorbing state phase transition, from this perspective, may be viewed as a probe into the stability of purely dissipative quantum state preparation.

The model consists of local Lindbladian dynamics that includes the generalized W state as a dark state. These are a class of states given by

$$|W_{\text{Gen}}\rangle = \frac{1}{\sqrt{N}} \sum_{a=0}^{N-1} e^{i\phi_a} |\downarrow_0 \dots \downarrow_{a-1} \uparrow_a \downarrow_{a+1} \dots \downarrow_{N-1}\rangle, \quad (1)$$

where a labels the sites of the lattice, and various choices of phase ϕ_a give, for example, the W state (all $\phi_a = 0$) or vortex lattices (phase winding around closed loops). While the (generalized) W state is a dark state of the Lindbladian (to

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be defined below), some terms are “attractive” in the sense that they drive the system toward the W state (the W -state preparation protocol) while other terms are “repulsive” and drive states in the neighborhood of the W state toward a different state (in this case, the maximally mixed state); see the end of the next section for a precise definition as well as the Supplemental Material [53] for further details. Competition between these terms induces a phase transition away from the absorbing W state.

II. THE MODEL

The model we consider is a lattice of two level systems where pairs of lattice sites are coupled (via a set of environments) to undergo dissipative evolution. We restrict to the case well described by Lindbladian dynamics, given by the master equation for the density operator ρ :

$$\dot{\rho} = \sum_{\ell} \gamma_{\ell} \left(L_{\ell} \rho L_{\ell}^{\dagger} - \frac{1}{2} \{ L_{\ell}^{\dagger} L_{\ell}, \rho \} \right). \quad (2)$$

The model is purely dissipative (with real $\gamma_{\ell} > 0$), which is why no coherent, Hamiltonian evolution appears in (2).

In particular, we consider two-site jump operators grouped into the following four sets:

$$L_1 = \{|1, 0\rangle_{a,b} \langle 0, 0|_{a,b}\}, \quad (3a)$$

$$L_2 = \{|1, 0\rangle_{a,b} \langle 1, 1|_{a,b}\}, \quad (3b)$$

$$L_3 = \{|0, 0\rangle_{a,b} \langle 1, 1|_{a,b}\}, \quad (3c)$$

$$L_{3'} = \{|1, 1\rangle_{a,b} \langle 0, 0|_{a,b}\}. \quad (3d)$$

We will consider two cases: all-to-all jump operators—with any pair of sites a, b —as well as nearest-neighbor (NN) where a, b are on adjacent lattice sites (with coordination number Z). We have written two-site states in terms of total spin quantum numbers $|s, m_s\rangle$.

In the Lindbladian (2), terms L_1, L_2, L_3 , and $L_{3'}$ have coefficients $\gamma_1, \gamma_2, \gamma_3$, and $\gamma_{3'}$, respectively. Throughout the work, we will take $\gamma_3 = \gamma_{3'}$. The model is summarized in Fig. 1(a).

To see that this Lindbladian represents a quantum reaction-diffusion model, view the spin-up excitations as hard-core bosonic particles. From this viewpoint, L_1 conserves particle number (total S_z) and—since the dissipative evolution favors states where a particle is in a coherent superposition on neighboring sites—will tend to spread the particle throughout the system, i.e., it is the quantum analog of the “diffusion” term. On the other hand, the rest of the jump operators correspond to prototypical reaction processes. Namely, “coagulation” terms L_2 and L_3 correspond to processes where two particles collide (i.e., are on adjacent sites) and decay down to a single particle state, while “branching” terms $L_{3'}$ correspond to Lindblad jumps from a single particle state to a two particle state. Two-site dissipative dynamics similar to that generated by the quantum jumps above have been realized across a variety of experimental platforms, for example [1, 14, 54].

From the quantum state preparation perspective, the model exhibits competition between quantum jumps where the W state is either an “attractive” (L_1, L_2) or “repulsive” ($L_3, L_{3'}$) dark state. These properties are defined by the structure of the

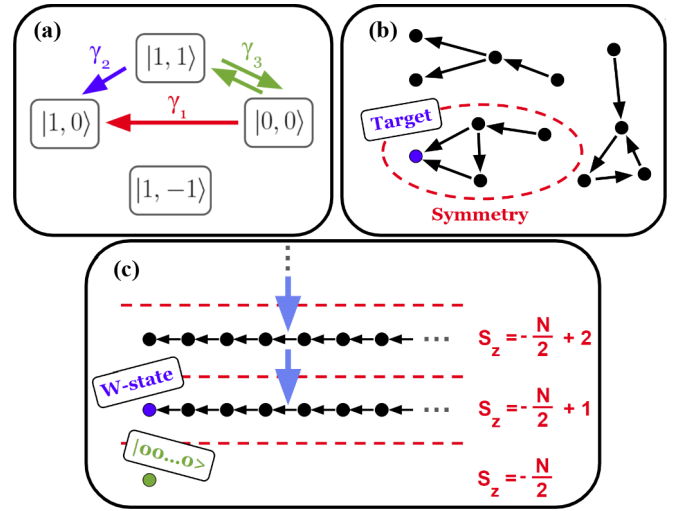


FIG. 1. Our model and dissipative W -state preparation. (a) Quantum jump operators for the model written in terms of total spin eigenstates $|s, m_s\rangle$ for two-site pairs. The γ_1, γ_2 terms prepare the W state, while γ_3 terms induce a phase transition to a mixed state. (b) Points represent states and arrows are quantum jumps in the Lindblad dynamics. (c) Preparation of W state using ladder between symmetry sectors.

adjoint of the jump operators, $L^{\dagger}$. Attractive dark states satisfy $L^{\dagger}|W\rangle \neq 0$; this implies there is some attractive basin of states which flow into the W state under L . For repulsive dark states, $L|W\rangle = L^{\dagger}|W\rangle = 0$. Here, states in the neighborhood of the W state will never jump to the W state (i.e., $\langle W|\rho(t)|W\rangle = \langle W|\rho(0)|W\rangle$), but will instead flow to some other steady state. See [53] for further details.

III. DISSIPATIVE W -STATE PREPARATION

We now show that Lindbladian dynamics with L_1 and L_2 prepares the (generalized) W state. While this is not feasible in full generality—the generalized W state falls outside the class of states preparable through local, purely dissipative Lindbladians [11]—various strategies have been explored to extend the set of preparable states [12, 55–58] allowing for experimental W -state realizations [59–61].

Here, the target state is prepared from a restricted subspace of allowed initial states. Jump operators which commute with a symmetry of the target state may be used for dissipative preparation within the subspace of states that share that symmetry—Fig. 1(b). Additional jump operators take other states into the target symmetry sector, broadening the space of initial states which prepare the target.

To prepare the W state, we focus on subspaces of states with a fixed total S_z . Within the one-spin-up ($S_z = -\frac{N}{2} + 1$) sector, the W state is the only state with maximal total spin. Applying the local jump operators L_1 —which preserve S_z while increasing total spin—therefore drives the one spin-up subspace to the W state. In other words, these jump operators locally lock the phase between sites, driving toward global phase coherence—an idea previously considered in [10, 62] for similar jump operators in lattice ultracold bosonic atoms.

Adding L_2 —or any L which takes a two spin-up state to a one spin-up state, including L_3 —connects higher- S_z sectors to the target $S_z = -\frac{N}{2} + 1$ sector. The full protocol, shown in Fig. 1(c), combines L_1 to push each S_z -sector toward globally phase coherent states and L_2 to reduce S_z until fewer than two spin ups remain. In this way, any initial state will converge to the W state except for one—the vacuum (i.e., completely spin-down state $|\downarrow \downarrow \dots \downarrow\rangle$). Like the W state, the vacuum is a dark state of L_1 and L_2 . Given the no-go theorem [11], this is the closest possible procedure to fully dissipative preparation of the W state; only one state is not driven by the dissipative attractor dynamics to the W state, so a random initial state will converge to the W state with probability one in the thermodynamic limit. In the Supplemental Material [53], we analytically show preparation time for this protocol scales as $O(N^2)$.

To prepare the generalized W state, note that the generalized W state is simply a local gauge transformation of the W state, i.e.,

$$|W_{\text{Gen}}\rangle = U_{\text{gauge}}|W\rangle = \prod_a e^{i\phi_a \sigma_a^+ \sigma_a^-} |W\rangle \quad (4)$$

with $\sigma_a^\pm$ the spin raising/lowering operators and ϕ_a desired phase to imprint on site a . A preparation strategy may then be constructed by applying this transformation to the jump operators, $L_\ell \rightarrow U_{\text{gauge}} L_\ell U_{\text{gauge}}^\dagger$, such that the dynamics under (2) is preserved except to the new absorbing state $|W_{\text{Gen}}\rangle$.

IV. ABSORBING STATE PHASE TRANSITION

We now investigate what happens when $\gamma_3 \neq 0$. Initially, we will focus on the all-to-all case for two main reasons.

First, the W state is all-to-all symmetric and so leveraging an all-to-all preparation procedure might be expected to improve upon preparation time. Indeed, the all-to-all model prepares the W state in constant time as opposed to the polynomial scaling of the NN case. This improvement, however, comes at a cost; the all-to-all procedure is fragile to $O(\frac{1}{N})$ error quantum jump rates, whereas the transition occurs for error jump rates $O(\frac{1}{Z})$ in the NN case.

Second, any all-to-all Lindbladian (or general open quantum dynamics) may be represented by a set of generalized Dicke operators, which act as closed operations within the space of generalized Dicke states [63]. Specifically, any permutation symmetric density matrix may be represented as a linear combination of generalized Dicke states,

$$\rho = \sum_{q,q_z,\sigma_z} \alpha_{q,q_z,\sigma_z} \mathcal{D}_{q,q_z,\sigma_z}, \quad (5)$$

where the generalized Dicke states $\mathcal{D}_{q,q_z,\sigma_z}$ are defined in terms of three quantum numbers q , q_z , and σ_z . The number of $\mathcal{D}_{q,q_z,\sigma_z}$ is only $O(N^3)$, and therefore scales favorably compared to the exponential size of the full density matrix. Generalized Dicke states extend to open systems, the more commonly used Dicke states, i.e., the set of permutation symmetric pure states whose utility was first introduced in the context of quantum optics [64]. We will focus here on the main results from the generalized Dicke analysis; calculation details may be found in the Supplemental Material [53].

Using generalized Dicke operators, it is possible to see explicitly that the only two dark states of the system with $\gamma_3 = 0$ are the vacuum and the W state. Furthermore, any initial state that has zero overlap with the vacuum will converge to the W state. The operators L_3 , $L_{3'}$ on the other hand have three important steady states: the vacuum, the W state, and the maximally mixed state (within the state space orthogonal to the two dark states). Importantly, as remarked in our introduction of the model, both the vacuum and W state are repulsive dark states. Initial states with zero overlap in the vacuum or W state will flow to the maximally mixed state.

To numerically study the phase transition induced by these competing jump operators, we examine the gap Δ to the first nonzero eigenvalue. The gap closes as γ_3 is increased (and $\gamma_1 = \gamma_2 = 1$). We plot the gap in Fig. 2(a) as a function of $N\gamma_3$. The gap closing for such small values of γ_3 —scaling with $\frac{1}{N}$ in the system size—implies that the W -state preparation strategy is fragile to the addition of L_3 , $L_{3'}$.

A priori, this fragility may be surprising as L_1 and $L_{3'}$ jump from the same state (the singlet state) and act on the same number of site-pairs (similarly for L_2 , L_3 , and the $|1, 1\rangle$ state). There is, therefore, no $O(N)$ difference in the rate at which cohering vs decohering quantum jumps occur. This would seem to suggest a transition should occur when $\frac{\gamma_3}{\gamma_1} = O(1)$ not $O(\frac{1}{N})$.

To see why this is not the case, it is helpful to look at the NN model and consider evolution of the expectation value $\langle P^{[1,0],[1,-1]} \rangle$ where $P^{[1,0],[1,-1]} = \frac{1}{N-1} \sum_{(a,b)} (|1, 0\rangle_{a,b} \langle 1, 0|_{a,b} + |1, -1\rangle_{a,b} \langle 1, -1|_{a,b})$. Physically, $\langle P^{[1,0],[1,-1]} \rangle$ represents the expectation value that neighboring sites will be in either the $|1, 0\rangle$ or $|1, -1\rangle$ state. This is maximal for the W state and serves as an order parameter for the transition. As an example, consider evolution of this expectation value for any ρ in the $S_z = -\frac{N}{2} + 1$ sector (see the Supplemental Material for derivation with any ρ [53]) given by

$$\frac{d}{dt} \langle P^{[1,0],[1,-1]} \rangle = \left(\gamma_1 - \gamma_3 \frac{Z-1}{2} \right) \langle P^{[0,0]} \rangle, \quad (6)$$

where $\langle P^{[0,0]} \rangle$ is the expectation value that neighboring sites are in the $|0, 0\rangle$ state. Intuitively, the extra factor of Z occurs for the decohering terms because—for any pair of sites a, b —the jump operators can decohere a, b by acting on a and any neighbor $c \neq b$ (or similarly, sites b and a neighbor $c \neq a$). However, to make a, b coherent, L_1 must act directly on sites a, b . In the all-to-all limit, $Z \rightarrow N$ and the decohering terms dominate by a factor of N , implying any nonzero γ_3 will ruin the W -state preparation procedure.

We now investigate the properties of the steady state in the regime where the gap closes. Here, the steady state will be a linear combination of the W -state and the eigendensity-matrix corresponding to the third zero eigenvalue. The relative weights of the linear combination may in general depend on the initial state and the values of $\gamma_1, \gamma_2, \gamma_3$. In Fig. 2(b), we compare steady states taking as our initial state the totally spin polarized system $|\uparrow \uparrow \dots \uparrow\rangle$. The steady state is a mixture of density matrices with well defined total S_z given by $\rho_{\text{steady}} = \sum_{S_z} \alpha_{S_z} \rho_{S_z}$ with $\sum_{S_z} \alpha_{S_z} = 1$. We plot α_{S_z} [given by

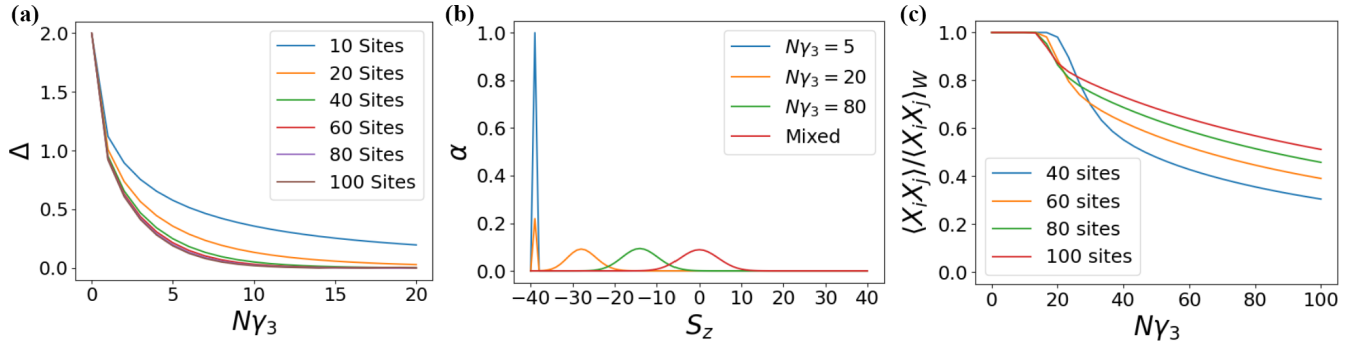


FIG. 2. Absorbing state phase transition in the all-to-all system. (a) Gap to first nonzero eigenvalue. Here, $\gamma_1 = \gamma_2 = 1$ and numerics are done within the $\sigma_z = 0$ subspace to access larger system sizes. (b) Proportion of Steady state (with totally spin-up initial state) in each S_z symmetry sector. Numerics were done with $N = 80$ and run for $t = 1000$. (c) Long range correlations in the steady state. Transition point is consistent with closing of the gap in (a).

$\alpha_{\frac{N}{2}, S_z, 0}$ in the generalized Dicke representation (5)] in Fig. 2(b) to show the relative weight of the steady state in each S_z sector.

When $\gamma_1, \gamma_2 \ll \gamma_3$, the steady state [labeled as “mixed” in Fig. 2(b)] is the maximally mixed state (minus the all spin down and W state). From counting the number of states in each total S_z sector, the distribution of coefficients for this “mixed” state is given by $\alpha_{S_z}^{\text{Mixed}} = \frac{\binom{N}{\frac{N}{2} + S_z} - \delta_{S_z, -\frac{N}{2} + 1}}{2^N - 2}$, where δ is the Kronecker delta and the -2 comes from the removal of the all spin down and W state. As γ_3 decreases, the distribution shifts from the maximally mixed state to a distribution centered around smaller and smaller S_z . Additionally, the weight of the W state increases as γ_3 becomes small (with $\alpha_{-\frac{N}{2} + 1}$ increasing toward one). In the regime where the gap opens, the steady state becomes the W state.

As a measure of coherence in the system, we consider the two-point correlator $\langle X_a X_b \rangle$ for Pauli X_a at site a . For the W state, we have $\langle X_a X_b \rangle_W = \frac{2}{N}$. In Fig. 2(c), we compare $\langle X_a X_b \rangle$ for the steady state of our model to that of the W state. The ratio $\frac{\langle X_a X_b \rangle}{\langle X_a X_b \rangle_W}$ departs from one at the same value of γ_3 where the gap closes [compare to Fig. 2(a)]. Coherence still persists, however, after the transition.

The steady state, therefore, does not just converge to the steady state of L_3 (the maximally mixed state) after the transition. Instead, the steady state is a mixture of states with coherence less than, but comparable to, that of the W state. It is only in the limit $\gamma_3 \gg \gamma_1, \gamma_2$ that this coherence approaches zero and the steady state becomes maximally mixed.

V. NN RESULTS AND CRITICALITY

We now contrast the above all-to-all results with the NN version of the model, as well as discuss critical behavior. As previously discussed, the transition in the NN model occurs when $\frac{\gamma_3}{\gamma_1} = O(1)$. Furthermore, unlike the all-to-all transition which goes from gapped to gapless, the NN transition is from a polynomial gap to an exponential gap in system size. The replacement of the fixed gap (for all-to-all) with a polynomial gap is because the local jump operators in the NN model need time at least on order of the system size to generate coherence across the entire system. In the mixed phase, the NN model also exhibits coherence, but in this case the coherence decays exponentially in m for $\langle X_a X_{a+m} \rangle$ again

due to locality. Numerics for small system sizes supporting these results are provided in the Supplemental Material [53].

To analyze behavior at the critical point, we will focus on the dynamical critical exponent z given by the gap $\Delta = O(N^{-z})$. For 1D systems in the DP universality class, $z \approx 1.5807$ [36]. However, for the NN model, we must have $z \geq 2$. This is because the gap in the $\gamma_3 = 0$ limit (i.e., in the ordered phase) already scales as $\Delta = O(N^{-2})$ [53], and the addition of γ_3 terms may only increase the dynamical exponent. Small system size numerics yield the critical value $z \approx 2.7$ at the transition into the disordered phase. Thus, the critical behavior of the NN model is inconsistent with DP universality. The all-to-all model, on the other hand, did not show numerical inconsistencies with DP [53].

VI. CONCLUSIONS

We have investigated a purely dissipative quantum reaction-diffusion model which contains competing terms that act to either cohere or decohere the system. The cohering terms may be viewed as a dissipative quantum state preparation protocol for the generalized W state, while the decohering terms act as errors to the preparation procedure. This transition—from a long range coherent absorbing state to a mixed phase with short range coherence—leads to critical behavior which falls outside the directed percolation universality class typical for absorbing state phase transitions. This distinct behavior is due to the dynamics of coherence spreading throughout the system at criticality. We, furthermore, showed that while the all-to-all model can quickly prepare the W state, the preparation protocol is fragile to errors: the critical error rate scales as the inverse of the system size (whereas the nearest neighbor model is robust to weak error rates). This leads to a tradeoff between nearest neighbor models which prepare slowly but are robust to errors, and all-to-all models which prepare quickly but are fragile to errors. Investigating to what extent all-to-all dissipative state preparation procedures more generally are fragile to errors is an interesting avenue for future work.

While we have shown that coherence spreading at the transition can lead to non-DP critical behavior, many open questions remain surrounding the nature of this transition. For example, we have established distinction from DP through

a lower bound on the dynamical critical exponent, but it would be interesting to establish accurate estimates of all critical exponents through numerical or analytical methods. Furthermore, the universality of such critical exponents is also unclear as—for any dissipative preparation procedure with an absorbing state transition—the dynamical critical exponent must be lower bounded by the preparation time without errors. This suggests the possibility of distinct critical behavior depending on the dissipative preparation protocol used.

It is also an open question how much of a role the vacuum plays in our system. While it is disconnected from the dynamics, regions which are indistinguishable locally from the vacuum must also be dark and therefore will exhibit slowed dynamics. Such dynamics is typical in kinetically constrained

models, e.g., [65,66]. A full analysis of how this physics affects z and the dynamics more broadly we leave to future work. On the other hand, it is also possible to isolate the W state as the sole dark state through the addition of coherent terms [12]. Though more rigorous analysis is needed, preliminary small system size numerical results when this coherent term is added are also consistent with critical $z > 2$ [53].

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