Simulating the real-time evolution of large quantum systems is a notoriously hard problem. On the one hand, due to the enormous dimension of the Hilbert space, which grows exponentially with the system size, diagonalizing the Hamiltonian is impossible in practice. On the other hand, the configurations contributing to the real-time path integral have complex weights, which prevents the application of the Monte Carlo method based on importance sampling. While this method often works extremely well for Euclidean time simulations of quantum systems in thermal equilibrium, it fails for real-time simulations, due to a severe sign or complex weight problem. A notable exception are gapped 1-d systems with small entanglement, for which the matrix product states underlying the density matrix renormalization group [1, 2] provide a good basis for simulating the real-time evolution, at least for moderate time intervals [3-9]. Also Euclidean time simulations may suffer from severe sign problems, for example, in fermionic systems away from half-filling or in the presence of frustrating interactions. Some sign problems even fall in the complexity class of NP-complete problems [10], which can be solved in polynomial time on a hypothetical "non-deterministic" computer, but not on an ordinary deterministic computer (unless NP would unexpectedly coincide with the complexity class P). This means that a general method for solving sign problems is unlikely to exist, and that these problems should thus be addressed on a case by case basis. In fact, several severe sign problems have been solved completely using the meron-cluster algorithm [11, 12] or the fermion bag approach [13, 15].

It is not surprising that classical computers have problems simulating quantum systems, in particular, in real time. The entanglement inherent in complex quantum phases is not easily representable, let alone computable, as classical information. For this reason, as early as 1982 Feynman proposed using specifically designed quantum devices to mimic quantum systems that are difficult to simulate classically [15]. Since the ground-breaking experimental realization of Bose-Einstein condensation [17, 18], the fields of atomic physics and quantum optics have undergone impressive development. The degree to which ultracold atomic systems can be engineered and controlled is truly remarkable, and Feynman's vision of quantum simulators is becoming a reality. For example, the bosonic Hubbard model has been implemented with exquisitely well-controlled ultracold atoms in an optical lattice [19], and several aspects of this quantum simulation have been verified by comparison with accurate quantum Monte Carlo simulations [20]. Digital [21] and analog [22] quantum simulators are widely discussed in atomic and condensed matter physics [23, 28], and more recently also in a particle physics context [29, 37].

While quantum simulators are gradually becoming available, they are still far from being universally applicable, and they are not yet precision instruments. Hence, simulating the real-time evolution of large quantum systems on classical computers remains an important challenge. Since isolated quantum systems tend to evolve into complicated entangled states such as those of Schrödinger's cat, it will in general be extremely difficult to compute them classically. In the real world, Schrödinger cat states usually do not arise, because quantum systems suffer from decoherence by coupling to their environment, and thus behave more classically. It should hence be easier to simulate quantum systems in the presence of an environment. Here we develop a method to simulate the real-time evolution of large quantum spin systems whose dynamics are entirely driven by measurements of the total spin \( \langle S_x + S_y \rangle^2 \) of pairs of spins \( \frac{1}{2} \) at adjacent positions \( x \) and \( y \). The measurements give rise to a dissipative coupling to the environment, which drives the system from an initial state to a new equilibrium. Remarkably, when one averages over the measurement results, the sign problem is eliminated and the dynamics can be addressed with an efficient cluster algorithm. This is the first time that the real-time evolution of a large strongly coupled quantum system can be simulated over arbitrarily long time intervals in any spatial dimension. The dissipative measurement process that drives the time-evolution may even be realizable in optical lattice experiments. The control of quantum systems by measurements is investigated in [38, 39], and dynamical phenomena in out-of-equilibrium quantum systems are
discussed in \cite{10, 50}. Measurements have also been suggested as a resource for quantum computation \cite{51, 54}.

In non-relativistic quantum mechanics, the path integral representation of measurement processes has been discussed in \cite{55, 56}.

Let us consider a general quantum system with a (possibly time-dependent) Hamiltonian, whose real-time evolution from \( t_k \) to \( t_{k+1} \) is described by the time-evolution operator \( U(t_{k+1}, t_k) = U(t_k, t_{k+1})^\dagger \). At time \( t_k \) \((k \in \{1, 2, \ldots, N\})\) we assume an observable \( O_k \) is measured and an eigenvalue \( o_k \) is obtained as the measurement result. The Hermitian operator \( \rho_{o_k} \) projects on the subspace of the Hilbert space spanned by the eigenvectors of \( O_k \) with eigenvalue \( o_k \). Starting from an initial density matrix \( \rho_0 = \sum_i p_i |i\rangle \langle i| \) (with \( 0 \leq p_i \leq 1 \), \( \sum_i p_i = 1 \)) at time \( t_0 \), the probability of reaching a final state \( |f\rangle \) at time \( t_f \), after a sequence of \( N \) measurements with results \( o_k \), is then given by \cite{57}

\[
p_{\rho_0,f} (o_1, o_2, \ldots, o_N) = \sum_i \langle i | U(t_0, t_1) P_{o_1} U(t_1, t_2) P_{o_2} \ldots P_{o_N} U(t_N, t_f) | f \rangle 
\]

\[
\langle f | U(t_f, t_N) P_{o_N} \ldots P_{o_2} U(t_2, t_1) P_{o_1} U(t_1, t_0) | i \rangle p_i. \quad (1)
\]

The matrix elements of both the time-evolution and the projection operators are in general complex, thus leading to a severe sign problem in Monte Carlo simulations. As we have argued above, classical measurements disentangle the quantum system, at least to some extent, and should thus alleviate the sign problem. For simplicity, we now consider quantum systems whose time-evolution is entirely driven by measurements, i.e. \( U(t_k, t_{k+1}) = \mathbb{1} \). By inserting complete sets of states \( \sum_{n_k} |n_k\rangle \langle n_k| = \mathbb{1} \) into the first factor and independently \( \sum_{n_k'} |n_k'\rangle \langle n_k'| = \mathbb{1} \) into the second factor in eq. \((1)\), between the times \( t_k \), one arrives at a real-time path integral along the Keldysh contour leading from \( t_0 \) to \( t_f \) and back \cite{58, 59}. In the doubled Hilbert space of states \( |n_k n_k'\rangle \), encompassing both pieces of the Keldysh contour,

\[
p_{\rho_0,f} (o_1, o_2, \ldots, o_N) = \sum_i \langle i | \left( P_{o_1} \otimes P_{o_1}^* \right) \left( P_{o_2} \otimes P_{o_2}^* \right) \ldots \left( P_{o_N} \otimes P_{o_N}^* \right) | f \rangle \langle f | 
\]

\[
\sum_i \prod_{n_k} \langle n_{k-1} n_{k-1} | P_{o_k} \otimes P_{o_k}^* | n_{k} n_{k}' \rangle. \quad (2)
\]

We use the notation \( \langle n_{k-1} n_{k-1} | P_{o_k} \otimes P_{o_k}^* | n_{k} n_{k}' \rangle = \langle n_{k-1} | P_{o_k} | n_{k-1} \rangle \langle n_{k} | P_{o_k}^* | n_{k}' \rangle^* \) and \( \langle n_{N} n_{N}' \rangle = \langle f | f \rangle \). We also consider the probability \( p_{\rho_0,f} \) of reaching the final state \( |f\rangle \) irrespective of the intermediate measurement results,

\[
p_{\rho_0,f} = \sum_{o_1} \sum_{o_2} \ldots \sum_{o_N} p_{\rho_0,f} (o_1, o_2, \ldots, o_N) 
\]

\[
= \sum_i \prod_{n_k} \langle n_{k-1} n_{k-1} | \tilde{P}_k | n_{k} n_{k}' \rangle, \quad (3)
\]

where \( \tilde{P}_k = \sum_{o_k} P_{o_k} \otimes P_{o_k}^* \) is obtained by summing over all possible measurement results \( o_k \) at time \( t_k \).

Besides the process of sporadic measurements, let us also consider quantum systems that are continuously monitored by their environment. This situation is characterized by a set of Lindblad operators \( \{ \mathcal{L}_{o_k} \} \) (related to Kraus operators \( \{ K_{o_k} \} \) that obey \( (1 - \varepsilon \gamma N) \mathbb{1} + \sum_{i,o_k} \mathcal{L}_{o_k} \mathcal{L}_{o_k}^\dagger = \mathbb{1} \). Here \( \gamma \) determines the probability of measurements per unit time, and \( k \in \{1, 2, \ldots, N\} \) labels the operators \( O_k \) (with eigenvalues \( o_k \)) that can induce quantum jumps at any moment in time. In the continuous time limit, \( \varepsilon \to 0 \), and in the absence of a Hamiltonian, the time-evolution of the density matrix is then determined by the Lindblad equation

\[
\dot{\rho} = \frac{1}{\varepsilon} \sum_{k,o_k} \left( \mathcal{L}_{o_k} \rho \mathcal{L}_{o_k}^\dagger - \frac{1}{2} \mathcal{L}_{o_k}^\dagger \mathcal{L}_{o_k} \rho - \frac{1}{2} \rho \mathcal{L}_{o_k}^\dagger \mathcal{L}_{o_k} \right)
\]

\[
= \gamma \sum_{k,o_k} \left( \rho \mathcal{P}_{o_k} - \rho \right). \quad (4)
\]

As a simple example, let us first consider two spins \( \frac{1}{2} \), \( \vec{S}_x \) and \( \vec{S}_y \), forming total spin \( S \) eigenstates \( |S \rangle \) (with 3-component \( S^3 \)) \: |11\rangle = |1\rangle \langle 1|, \langle 00\rangle = \frac{1}{\sqrt{2}}(\langle \uparrow \downarrow + \downarrow \uparrow \rangle). \) The projection operators corresponding to a measurement 1 or 0 of the total spin are then given by \( P_1 = |11\rangle \langle 11| + 10 \langle 10| + 1 - 1 \langle 1 - 1| \) and \( P_0 = |00\rangle \langle 00| \), such that

\[
P_1 = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \frac{1}{2} & \frac{1}{2} & 0 \\ 0 & \frac{1}{2} & \frac{1}{2} & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}, \quad P_0 = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & \frac{1}{2} & -\frac{1}{2} & 0 \\ 0 & -\frac{1}{2} & \frac{1}{2} & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}. \quad (5)
\]

The negative entries in \( P_0 \) give rise to a sign problem in the corresponding real-time path integral. We quantize the spins in the 3-direction, with \( s_x = \pm \frac{1}{2} \) denoting the eigenvalues of \( S^3_x \). In the doubled Hilbert space of states \( |n_k n_k' \rangle = \sum_{s_x,k} |s_x,k's_y,k's_y,k' \rangle \) one then obtains

\[
\langle s_x,k's_y,k' | \mathcal{P}_1 | s_x,k+1's_y,k+1 \rangle = \delta_{s_x,k+1,k} \delta_{s_y,k+1,k} \delta_{s_x,k',k+1} \delta_{s_y,k',k+1} \delta_{s_y,k',k+1} + \delta_{s_x,k,k+1} \delta_{s_y,k',k} \delta_{s_x,k',k+1} \delta_{s_y,k',k+1} \delta_{s_y,k',k+1}/2. \quad (6)
\]

All matrix elements of \( \tilde{P} = P_1 \otimes P_1 + P_0 \otimes P_0 \) are non-negative. The Kronecker \( \delta \)-functions encode loop-cluster rules for binding parallel spins together \cite{63, 64}. The Lindblad process is the continuous-time limit of the discrete measurement process, which can be simulated directly in continuous time \cite{65}. Remarkably, the resulting cluster algorithm allows very efficient real-time simulations, without encountering a sign problem.

The simple two-spin system is easily extended to a large system in any dimension. We investigate a system of quantum spins \( \frac{1}{2} \) on a square lattice of size \( L \times L \) with periodic boundary conditions. To define an initial density matrix \( \rho_0 = \exp(-\beta H) \), we consider the antiferromagnetic Heisenberg Hamiltonian, \( H = J \sum_{(xy)} \vec{S}_x \cdot \vec{S}_y \),
which is used only to prepare an ensemble of initial states, not to evolve it further in time. The real-time evolution is again driven entirely by the measurement of the total spin \((\tilde{S}_x + \tilde{S}_y)^2\) of nearest-neighbor spin pairs. In a first step, all pairs of neighboring spins separated in the 1-direction, at \(x = (x_1, x_2)\) and \(y = (x_1 + 1, x_2)\) with even \(x_1\), are measured simultaneously. In a second step, the pairs at \((x_1, x_2)\) and \((x_1, x_2 + 1)\) with even \(x_2\), which are separated in the 2-direction, are examined. In a third and fourth measurement step, the total spins of pairs with odd \(x_1\) and \(x_2\) are being measured. Then the same four-step measurement process is repeated an arbitrary number of times \(M\), such that the total number of measurements is \(N = 4M\). This particular measurement sequence was chosen arbitrarily and can be replaced by any other one.

For the Lindblad process the ordering of the measurements is irrelevant. Together with the Keldysh contour, the Euclidean time interval \([0, \beta]\) (where \(T = 1/\beta\) is the temperature) forms a closed contour in the complex time plane. The clusters, which are closed loops extending through both real and Euclidean time, are updated by simultaneously flipping all spins belonging to the same cluster with probability \(\frac{1}{2}\). Remarkably, as a global consequence of eq. (6), the clusters, and thus also the spin configurations that contribute to the real-time path integral, are identical on both parts of the Keldysh contour, i.e., \(s_{x,k} = s'_{x,k}\). Then eq. (6) simplifies to

\[
\langle s_{x,k} s_{y,k} s_{x,k} s_{y,k} \rangle = \frac{1}{2} (\delta_{s_{x,k}, s_{y,k}} + \delta_{s_{x,k}, s_{y,k}+1} + \delta_{s_{x,k}, s_{y,k}+1} \delta_{s_{x,k}, s_{y,k}})/2 = \langle s_{x,k} s_{y,k} \rangle P_1 |s_{x,k}+1 s_{y,k+1}\rangle.
\]

We have investigated the real-time evolution of initial state ensembles corresponding to the 2-d square lattice Heisenberg antiferromagnet. While the uniform magnetization, \(M = \sum_x \tilde{S}_x\), i.e., the total spin, is conserved in the measurement process, the staggered magnetization, \(M'_0 = \sum_x (-1)^{x_1+x_2} \tilde{S}_x^3\), as well as the other Fourier modes \(\tilde{S}(p) = \sum_x \tilde{S}_x^3 \exp(ipx_1+ipx_2)\), \(p = (p_1, p_2)\), are affected by the measurements. Fig. 1a shows the staggered magnetization squared, averaged over the ensemble of final states \(|f\rangle\) that results after \(N\) discrete measurements, for systems with different initial temperatures. They are quickly driven to a new equilibrium ensemble.

In order to study the equilibration process in more detail, we now consider continuous Lindblad evolution, from an initial ensemble at low temperature \(\beta J = 5L/2a\), where \(a\) is the lattice spacing. Fig. 1c shows the real-time evolution of the Fourier modes

\[
|\langle \tilde{S}(p) \rangle^2 | \rightarrow A(p) + B(p) \exp(-t/\tau(p)),
\]

for a variety of momenta \(p = (p_1, p_2)\). While the conserved magnetization \(M\) with momentum \(p = (0, 0)\) does not equilibrate at all, low momentum modes equilibrate more slowly than high momentum modes. After a short initial phase, the various modes approach the ultimate new equilibrium exponentially, with an equilibration time \(\tau(p)\). Interestingly, for fixed momentum, \(\tau(p)\) is almost independent of the spatial volume. Large systems equilibrate slowly, because they contain modes of low momentum. For small momenta, the equilibration time behaves as \(1/|\gamma\tau(p)| = C |pa|^r, C = 1.26(8), r = 1.9(2)\) (Fig. 1b).

Since the measurement process conserves the total spin \(\tilde{S} = \sum_x \tilde{S}_x\), it does not change the probability distribution of the spin associated with the initial density matrix \(\rho_0\). The continuous-time Lindblad process even respects the translation and rotation symmetries of the lattice. The final density matrix, to which the system is driven by the measurements, is constrained by these symmetries, and is proportional to the unit matrix in each symmetry sector. This finally leads to a vanishing correlation length and to \(A(p) = L^4/4(L^2-1)\), indicated by the horizontal line in Fig. 1c. The unit density matrix (restricted to the appropriate symmetry sectors) is a stable \(T = \infty\) fixed point of any Hamiltonian plus Lindbladian dynamics, and thus a universal attractor for the ultimate long-term evolution for a large class of dissipative processes.

For the initial density matrix \(\rho_0\) of the antiferromagnet, \(\langle M_z^2 \rangle/L^2\) is proportional to \(L^2\), indicating spontaneous symmetry breaking of the \(SU(2)\) spin symmetry at zero temperature. By the Lindblad process the system is driven to a final density matrix for which \(\langle M_z^2 \rangle/L^2\)
becomes volume-independent, indicating that the $SU(2)$ spin symmetry is then restored. Consequently, the system must undergo a phase transition. Since the dissipative Lindblad process drives the system far out of thermal equilibrium, this phase transition is not expected to fall to any of the standard dynamical universality classes \[60\]. Figs. 2a,b show $\langle M_s^2 \rangle/L^4$ and the Binder ratio $\langle M_c^4 \rangle/\langle M_c^2 \rangle^2$ for $\beta J = 2L/3a$. The various finite-volume curves for the Binder ratio do not intersect. Instead, with increasing volume their inflection point moves to later times. Figs. 2c,d show the staggered magnetization density $M_s$ and the length scale $\xi = c/(2\pi \rho_s)$, where $c$ is the spinwave velocity and $\rho_s$ is the spin stiffness, as functions of time, obtained by a fit to

$$
\langle M_s(t)^2 \rangle = \frac{M_s(0)^2 L^4}{3} \sum_{n=0}^{3} c_n \left( \frac{\xi(t)}{L} \right)^n ,
$$

which implicitly defines $M_s(t)$ and $\xi(t)$. Here the constants $c_0 = 1$, $c_1 = 5.7503(6)$, $c_2 = 16.31(2)$, $c_3 = -84.8(2)$ (which are accurately determined at $t = 0$) are assumed to be time-independent. The order parameter $M_s(t) = M_s(0) \exp(-t/\tau)$ (with $M_s(0) = 0.30743(1)/a^2$) decays exponentially with $\gamma = 0.240(2)$, which suggests that the phase transition is completed only after an infinite amount of time. The length scale $\xi(t)$ (with $\xi(0) = 1.459(3)a$) increases with time, which can be attributed to a decrease of $\rho_s$.

By averaging over all measurement results, we have eliminated the sign problem. When one distinguishes individual measurement results, one encounters a sign problem. When one measures spin $S = 1$, one obtains

$$
\begin{align*}
\langle s_{x,k}s_{y,k}s_{x,k+1}s_{y,k+1} \rangle &= \langle \delta_{s_{x,k},s_{x,k+1}} \delta_{s_{y,k},s_{y,k+1}} + \delta_{s_{x,k},s_{y,k+1}} \delta_{s_{y,k},s_{x,k+1}} + \delta_{s_{x,k+1},s_{x,k}} \delta_{s_{y,k+1},s_{y,k}} + \delta_{s_{x,k+1},s_{y,k+1}} \delta_{s_{y,k+1},s_{x,k}} \rangle / 4 , \quad (10)
\end{align*}
$$

which is always non-negative. Again, the Kronecker $\delta$-functions encode rules for forming clusters of parallel spins. The four contributions to the right-hand side of eq. (10) correspond to four different cluster break-ups of the eight contributing spins. When we measure the total spin $S = 0$, we obtain

$$
\begin{align*}
\langle s_{x,k}s_{y,k}s_{x,k+1}s_{y,k+1} \rangle &= \langle \delta_{s_{x,k},s_{x,k+1}} \delta_{s_{y,k},s_{y,k+1}} + \delta_{s_{x,k},s_{y,k+1}} \delta_{s_{y,k},s_{x,k+1}} + \delta_{s_{x,k+1},s_{x,k}} \delta_{s_{y,k+1},s_{y,k}} + \delta_{s_{x,k+1},s_{y,k+1}} \delta_{s_{y,k+1},s_{x,k}} \rangle / 4 , 
\end{align*}
$$

FIG. 2. [Color online] a) $\langle M_s^2 \rangle/L^4$ and b) Binder ratio $\langle M_c^4 \rangle/\langle M_c^2 \rangle^2$ as functions of time for $L/a = 12, \ldots, 48$, $\beta J = 8, \ldots, 30$. Evolution of c) $M_s(t)/M_s(0)$ and d) $\xi(t)/\xi(0)$.
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