be modeled realistically. We use a hard-sphere model where each Rydberg excitation is described by a superatom [41] with a given blockade radius and effective Rabi frequency [21]. For more details on the parameters of the models, see the Appendix. We simulate the exact time-evolution of the experiment using the discrete truncated Wigner approximation (DTWA) [42]. DTWA simulations agree well with the experimental data as shown in Figs. 2(a)–2(c). The small deviation between simulations and experiments can be mostly attributed to an inaccuracy of the atom distribution obtained from the simplified excitation model (see Appendix).

The dynamics under the three different spin model in Figs. 2(a)–2(c) look strikingly similar in a log-linear plot. Indeed, by rescaling time with the characteristic timescale of each system given by $|J^{\perp}_{\text{median}} - J^{\parallel}_{\text{median}}|$, all relaxation curves coincide within the experimental errors. Here,

$$J^{\perp,\parallel}_{\text{median}} = \text{median}_j \max_i |J^{\perp,\parallel}_{ij}| \tag{2}$$

is the median of the nearest neighbor interaction strengths. This choice of typical interaction timescale is motivated by the oscillation frequency of a single pair of interacting spins governed by (1), which will be further discussed in the following section. The striking collapse allows us to infer the functional form of the relaxation dynamics of the nonintegrable models: For the Ising model, it is known that the relaxation follows exactly the stretched exponential law $e^{-(t/\tau)^{\beta}}$ [27] with stretching exponent β and timescale τ. The logarithm of the stretched exponential law is a power law. Plotted on a double logarithmic scale, this power law becomes a linear function [dashed line in the inset of Fig. 2(d)]. In this representation, the rescaled experimental data also show a linear behavior. This confirms the hypothesis that the stretched exponential law is the unifying description of the magnetization relaxation for the integrable quantum Ising model as well as the nonintegrable XX and XXZ Hamiltonians in the strongly disordered regime. We note that the dynamics are only robust with respect to a parameter of the microscopic model, the anisotropy $J^{\parallel}/J^{\perp}$, whereas the macroscopic geometry and also the dimension of the cloud may lead to different dynamics (see Appendix). In addition, we also measured the relaxation dynamics for various initial states (for one Hamiltonian) possessing different magnetization and again find similar relaxation dynamics at late times (see Appendix E).

IV. APPROXIMATE DESCRIPTION THROUGH STRONGLY INTERACTING PAIRS

In order to understand the regime where we have observed robust relaxation dynamics, we aim for a simplified model that includes only the relevant timescales of the system. To identify these, we exploit the strongly disordered nature of the system by adopting a perturbative approach in the spirit of the strong disorder renormalization group (SDRG) where the strongest coupling is integrated out iteratively [43–46].

In our model, the strongest coupled spins define a pair of spins. Crucially, the coupling within the pair will be much larger than all other couplings affecting the pair. This allows one to treat the coupling between this pair and the rest of the system perturbatively. To zeroth order, this pair of spins just decouples from the system and evolves independently. This elimination step, where we remove the strongest coupling, can be repeated within the rest of the system. For our initial state, each individual pair undergoes coherent dynamics between the fully polarized state in plus and minus x direction [see Fig. 3(a)] [47]. The resulting oscillation of the magnetization [shown in Fig. 3(b)] is independent of the specific XXZ Hamiltonian. Only the frequency, given by $J^{\perp}_{ij} - J^{\parallel}_{ij}$, differs depending on the Ising and exchange interaction strengths. This independence is at the origin of the observed model independence of relaxation dynamics.

With this model in hand, we can compute the time evolution of the magnetization by a simple average of cosine oscillations as shown by the grey dash-dotted lines (pair, noninteracting) in Figs. 3(c)–3(e). The resulting relaxation dynamics show good agreement with the experimental data. However, especially for the Ising and XXZ model, this model underestimates the timescales of the dynamics. This is somewhat expected, considering that the pair couplings found by iterative elimination are, on average, smaller than the nearest neighbor couplings.

Taking the perturbative treatment to next order, one finds an effective Ising-like coupling between pairs, as derived recently in the Appendix of [48]. The effective Hamiltonian governing the dynamics was found to be

$$\begin{aligned}\hat{H}_{\text{eff}} \approx \sum_{\langle i,j\rangle} \big(J^{\perp}_{ij}/2(\hat{s}^i_+\hat{s}^j_- + \hat{s}^i_-\hat{s}^j_+) + J^{\parallel}_{ij}\hat{s}^i_z\hat{s}^j_z\big). \\ + \sum_{\langle i,j\rangle,\langle k,l\rangle} J^{\text{eff}}_{ijkl}\hat{s}^{(i)(j)}_z\hat{s}^{(k)(l)}_z \end{aligned} \tag{3}$$

$$J^{\text{eff}}_{ijkl} = J^{\parallel}_{ik} + J^{\parallel}_{il} + J^{\parallel}_{jk} + J^{\parallel}_{jl} \tag{4}$$

where $\langle i,j\rangle$ denotes the summation over paired spins i and j and $2\hat{s}^{(i)(j)}_z = \hat{s}^i_z + \hat{s}^j_z$.

Fortunately, this model is integrable and allows for derivation of an analytical solution for the evolution of $\langle\hat{S}_x(t)\rangle$ (see Appendix F), which reads

$$\langle\hat{S}^{\text{pair}}_x\rangle(t) = \frac{1}{N}\sum_{\langle i,j\rangle}\cos\left(\frac{1}{2}(J^{\perp}_{ij} - J^{\parallel}_{ij})t\right)\prod_{\langle k,l\rangle}\cos^2\left(\frac{1}{8}J^{\text{eff}}_{ijkl}t\right). \tag{5}$$

The first factor in each term originates from the pair dynamic to zeroth order, as described previously. The other factors are reminiscent of the Emch-Radin solution for the Ising model and stem from the effective Ising interaction among the pairs. This effective Ising model of pairs captures the overall demagnetization dynamics remarkably well for all observed times [see Figs. 3(c)–3(e)], yielding very similar (and in the case of XXZ, even better) results compared to dTWA.

From the analytical form of the time evolution, Eq. (5), we find that many different oscillation frequencies contribute to each spin's magnetization dynamics. Most of these frequencies are very small, however, and do not contribute to the early-time dynamics. Thus, a reasonable ansatz for rescaling to make the dynamics collapse is to consider only the fastest frequency for each spin. Due to the highly disordered nature of our system, this strongest coupling will essentially always correspond to the closest neighboring spin.