

Nondestructive Optical Readout and Manipulation of Circular Rydberg Atoms

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(Received 24 October 2025; accepted 3 April 2026; published 20 May 2026)

Among the thriving quantum computation and quantum simulation platforms based on arrays of Rydberg atoms, those using circular Rydberg atoms are particularly promising. These atoms uniquely combine the strong dipole-dipole interactions typical of Rydberg states with long lifetimes. However, low-angular-momentum (ℓ) laser-accessible Rydberg levels have been so far mostly used, because circular Rydberg atoms have no optical transitions, hindering their individual detection and manipulation. We remove this limitation with a hybrid platform, combining an array of data laser-trapped circular Rydberg atoms of rubidium with an auxiliary array of Rb ancilla atoms transiently excited to a low- ℓ Rydberg level. We perform a quantum nondemolition detection of the data qubit with the ancilla, through the blockade of the ancilla optical excitation induced by a Förster resonance. Conversely, we locally manipulate the data qubit through the excitation of the ancilla. This dual-Rydberg platform is highly promising for quantum computation and simulation. It adds to the circular-atom toolbox the midcircuit measurements, essential for error correction. More strikingly, it gives access to time correlations in long-term quantum simulations, uniquely accessible to circular Rydberg atoms.

DOI: [10.1103/st3l-14d7](https://doi.org/10.1103/st3l-14d7)

Subject Areas: Atomic and Molecular Physics,
Quantum Physics

Rydberg atoms, i.e., atoms with one valence electron excited to a high-principal-quantum-number level, have recently proven to be ideal for quantum computation and quantum simulation [1,2]. The devices typically operate on arbitrary defect-free arrays of individual ground-state atoms laser trapped in optical tweezers [3,4]. The atoms are transiently excited to Rydberg levels. Their strong mutual dipole-dipole interaction enables two-qubit or multiqubit gates [5,6] and quantum simulation of interacting-spin Hamiltonians [7,8]. Remarkable recent achievements include the demonstration of high-fidelity gates [9–11], the operation of a quantum processor [12], or the quantum simulation of dynamics of interacting spins [13–15], with performances ultimately limited by the $\approx 100\text{-}\mu\text{s}$ lifetime of the laser-accessible Rydberg levels used in these devices. The final readout of the atomic state relies on fluorescence imaging of a ground state. While it can reach high fidelities and can be made fast [16], it is necessarily destructive for a qubit encoded in Rydberg levels.

Hybrid platforms, where ancillae are used to manipulate and detect the data qubits, are promising to solve this problem. Dual-species experiments propose to encode data qubits and ancillae in different atomic or molecular species [17–20], while coupling them strongly via the excitation to Rydberg levels experiencing a Förster resonance [21,22]. This enables, in particular, final or midcircuit nondestructive quantum state measurements of the data qubit [17] and multiqubit gates [23]. Recent experimental demonstrations of this architecture employing alkali atoms [24,25] have, in particular, enabled midcircuit correction of errors [26]. These methods require simultaneous advanced manipulation of two different species, increasing significantly the experimental complexity.

In this work, we employ different Rydberg levels of the same atomic species, namely rubidium 87 (Rb). We encode the data qubit on two long-lived circular Rydberg levels [27], while we engineer the ancilla qubit with two hyperfine sublevels of the ground state. We perform an ancilla-based measurement (ABM) of the state of the data atom, with a measurement fidelity $\mathcal{F} = 96\%$, using the dipole-dipole interaction with the ancilla, optically excited to a low-angular-momentum (ℓ) Rydberg level [28]. It possesses a Stark-tuned Förster resonance with one of the circular Rydberg levels. Because the circular states are impervious to optical light but highly sensitive to microwave (MW) fields, the two qubits can be separately manipulated with global addressing fields, hence acting as

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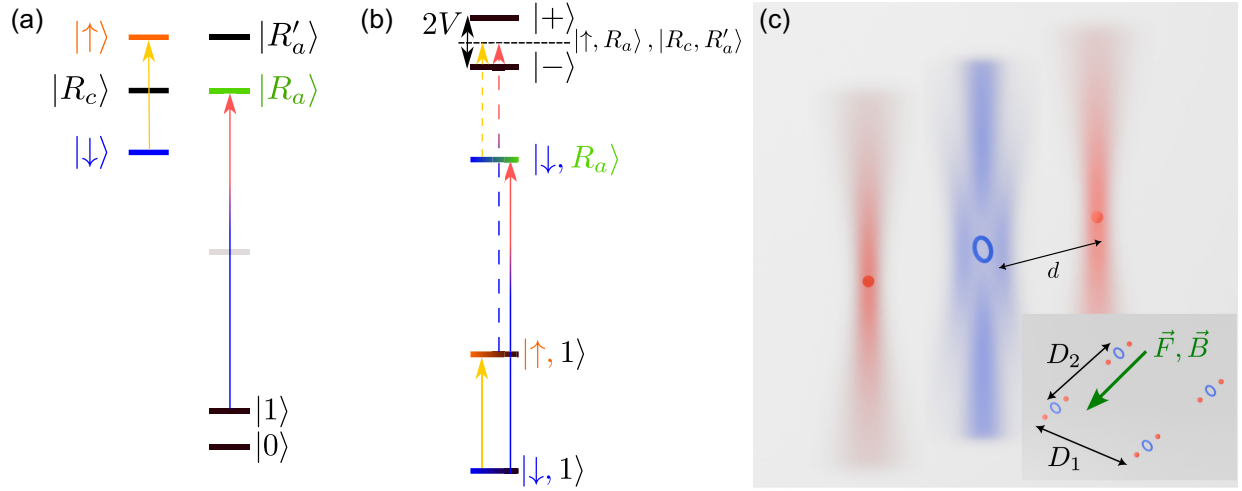


FIG. 1. Level structure and experimental geometry. (a) Simplified structure of noninteracting Rb data and ancilla atomic levels, not to scale. The data atom (left) is initially excited to the circular Rydberg level with principal quantum number $n = 52$, $|\downarrow\rangle$. A two-photon MW transition (yellow arrow) couples it to $|\uparrow\rangle$ ($n = 54$ circular state). The ancilla (right) is excited from its ground state $|1\rangle \equiv |5S_{1/2}, F = 2, m_F = 2\rangle$ to the Rydberg level $|R_a\rangle \equiv |45S_{1/2}, m_J = 1/2\rangle$ by a two-photon laser excitation (blue-red gradient arrows) through the $6P_{3/2}$ intermediate level (gray line). The state $|0\rangle$ is any ground state in the $|5S_{1/2}, F = 1\rangle$ manifold. At the Förster resonance, the energy difference between $|R'_a\rangle \equiv |45P_{3/2}, m_J = 3/2\rangle$ and $|R_a\rangle$ matches that between $|R_c\rangle$ ($n = 53$ circular state) and $|\uparrow\rangle$. (b) Simplified level structure of a pair of interacting data and ancilla atoms. Both optical and microwave excitations are allowed (solid arrows) from the pair state $|\downarrow, 1\rangle$ but blocked (dashed arrows) from $|\uparrow, 1\rangle$ and $|\downarrow, R_a\rangle$ because of the dipole-dipole interaction V that lifts the degeneracy between the $|\pm\rangle$ states. The dashed line indicates the energy of the noninteracting pair states $|\uparrow, R_a\rangle$ and $|R_c, R'_a\rangle$. (c) An ensemble of three atoms, with a data atom excited to the circular Rydberg level $|\downarrow\rangle$ (blue torus) and two ancillae in the ground state $|1\rangle$ (red spheres). The ancillae are trapped in Gaussian optical tweezers (red beams) and the data atom in an optical bottle beam (blue beam). The programmed interatomic distance is $d = 6 \mu\text{m}$. Inset: we prepare four ensembles distant by $D_1 = 50 \mu\text{m}$ and $D_2 = 60 \mu\text{m}$. Within an ensemble, the atoms are aligned along the quantization axis, fixed by the parallel electric $\vec{F}$ and magnetic $\vec{B}$ fields (green arrow).

atoms from different species. We demonstrate that the ABM is a quantum nondemolition (QND) measurement [28–30]: It projects the state of the data atom onto the measured state, according to Born probabilities, and leaves an atom initially in an eigenstate of the measurement unaffected. In addition, we use the interaction between the ancilla and data atoms for local manipulation of the state of the latter. Through local optical control of the ancilla atom only, we induce a phase flip of the data qubit with a fidelity of 93%, and locally induce a state transfer with a selectivity of 95%.

We encode the states of the data qubits onto the circular Rydberg levels with principal quantum numbers $n = 52$ and $n = 54$, denoted $|\downarrow\rangle$ and $|\uparrow\rangle$, respectively [see Fig. 1(a)]. The ancilla qubit states, denoted $|1\rangle$ and $|0\rangle$, correspond to the $|5S_{1/2}, F = 2, m_F = 2\rangle$ ground state and to any state of the $F = 1$ hyperfine ground-state manifold, respectively. We use a microwave field to induce transitions between $|\downarrow\rangle$ and $|\uparrow\rangle$ and laser beams to couple $|1\rangle$ to the Rydberg level $|R_a\rangle \equiv |45S_{1/2}, m_J = 1/2\rangle$. The latter resonantly interacts with $|\uparrow\rangle$ only, through a Förster resonance. In an electric field of $F = 1.6 \text{ V cm}^{-1}$, the atom-pair state $|\uparrow, R_a\rangle$ is degenerate with $|R_c, R'_a\rangle$, where $|R'_a\rangle = |45P_{3/2}, m_J = 3/2\rangle$ and $|R_c\rangle$ is the $n = 53$ circular state (see Appendix B). The electric dipole-dipole

interaction couples the states $|\uparrow, R_a\rangle$ and $|R_c, R'_a\rangle$ and lifts their degeneracy: The two new eigenstates $|\pm\rangle = (|\uparrow, R_a\rangle \pm |R_c, R'_a\rangle)/\sqrt{2}$ are displaced by $\pm V$ [see Fig. 1(b)], where the dipole-dipole interaction strength V scales with the interatomic distance d as $1/d^3$. At first order, the state $|\downarrow, R_a\rangle$ is unaffected by the dipole-dipole interaction. The strong first-order dipole-dipole interaction between $|\uparrow\rangle$ and $|R_a\rangle$ results in the Rydberg blockade [31] of the optical excitation from $|1\rangle$ to $|R_a\rangle$ of an ancilla atom when a nearby data atom is prepared in $|\uparrow\rangle$ [Fig. 1(b)]. This excitation is possible when the data qubit is in $|\downarrow\rangle$. Conversely, an ancilla in $|R_a\rangle$ prevents the microwave transfer between $|\downarrow\rangle$ and $|\uparrow\rangle$ for the data atom. We use these two processes to demonstrate QND measurement and local manipulation, respectively, of the state of the data qubit.

The experimental setup is described in detail in Ref. [32] and in Appendix A. We initially prepare a defect-free array of Rb atoms, laser trapped in focused Gaussian optical tweezers. We arrange the atoms in ensembles of three trapping sites, separated by a distance set to be $d = 6 \mu\text{m}$, and aligned with the quantization axis [see Fig. 1(c)]. This axis is fixed by a constant magnetic field $\vec{B}$ and a parallel initially vanishing electric field $\vec{F}$. The ensembles are far enough from each other to neglect interactions between atoms of different ensembles.

In every ensemble, we use an atom in the central site [Fig. 1(c)] as a data atom while atoms in the two outer sites act as ancillae. We can choose to operate with one or two ancillae, depending on the experiment we perform. We denote $|\psi_D, \psi_A\rangle_{DA}$ ($|\psi'_A, \psi_D, \psi_A\rangle_{ADA}$) the state of an atom pair (triplet) with the data atom and ancilla(e) in $|\psi_D\rangle$ and $|\psi_A\rangle$ or $|\psi'_A\rangle$, respectively. All the atoms are initially prepared in $|1\rangle$ and held in Gaussian optical tweezers.

For the time being, we consider for simplicity the situation with only one ancilla. The preparation of $|\downarrow, 1\rangle_{DA}$ begins with the laser excitation of the data atom to the low- ℓ Rydberg level $|52D_{5/2}, m_J = 5/2\rangle$. During this step, we turn off the trapping light for the central atom only, so that the ancilla remains in $|1\rangle$ because of the trap-induced light shifts. Then, the data atom is transferred to $|\downarrow\rangle$ using a circularization method described in Refs. [27,32]. This process leaves $|1\rangle$ unaltered. A final optional MW π pulse on the two-photon $|\downarrow\rangle \rightarrow |\uparrow\rangle$ transition prepares $|\uparrow, 1\rangle_{DA}$. The data atoms, in circular Rydberg states, are trapped in optical bottle beams (BoBs) [33] [Fig. 1(c)] that we turn on right after their initial optical excitation [32], while the ancillae remain trapped in the Gaussian optical tweezers.

At the end of the experimental sequence, we optically detect the state of the circular Rydberg atoms with a destructive, nonrepeatable measurement (DM) demonstrated in Ref. [32]. We revert the excitation process to bring the data atoms in $|\downarrow\rangle$ back to $|1\rangle$ and recapture them in Gaussian optical tweezers. All spuriously populated levels and $|\uparrow\rangle$ end up in Rydberg states and are expelled by the optical tweezer traps. A final 25-ms-long fluorescence imaging detects laser-trapped ground-state atoms, measuring the population in $|\downarrow\rangle$ prior to the DM. We can measure instead the population in $|\uparrow\rangle$ by applying an MW π pulse on the $|\downarrow\rangle \rightarrow |\uparrow\rangle$ transition before this measurement sequence. In both cases, we measure the probability to recapture the data atom in the optical tweezers, and denote it $\mathcal{P}_D^{(\uparrow)}$ or $\mathcal{P}_D^{(\downarrow)}$, the exponent indicating the detected state.

The *ABM sequence* is sketched in Fig. 2(a). (i) We turn off the traps and (ii) shine the Rydberg-excitation lasers on resonance with the $|1\rangle \rightarrow |R_a\rangle$ transition for 0.4 μ s. If the data atom is in $|\psi_D\rangle = |\downarrow\rangle$, the optical π pulse brings $|\downarrow, 1\rangle_{DA}$ to $|\downarrow, R_a\rangle_{DA}$. If $|\psi_D\rangle = |\uparrow\rangle$, however, Rydberg blockade prevents the excitation [see Fig. 1(b)], leaving $|\uparrow, 1\rangle_{DA}$ unaffected. At this stage, the state of the data atom, $|\downarrow\rangle$ or $|\uparrow\rangle$, is mapped onto the state of the ancilla, $|R_a\rangle$ or $|1\rangle$. (iii) We shelve the atoms left in $|1\rangle$ into $|0\rangle$ by optical pumping and (iv) we transfer the atoms excited to $|R_a\rangle$ back to $|1\rangle$ with the same optical π pulse as in step (ii), avoiding limitations due to the 51- μ s finite lifetime of $|R_a\rangle$. (v) We expel these atoms in $|1\rangle$ from the trapping region with a push-out optical beam, (vi) we turn back on the Gaussian optical tweezers to recapture the atoms left in $|0\rangle$, and we eventually detect all ground-state atoms by a final

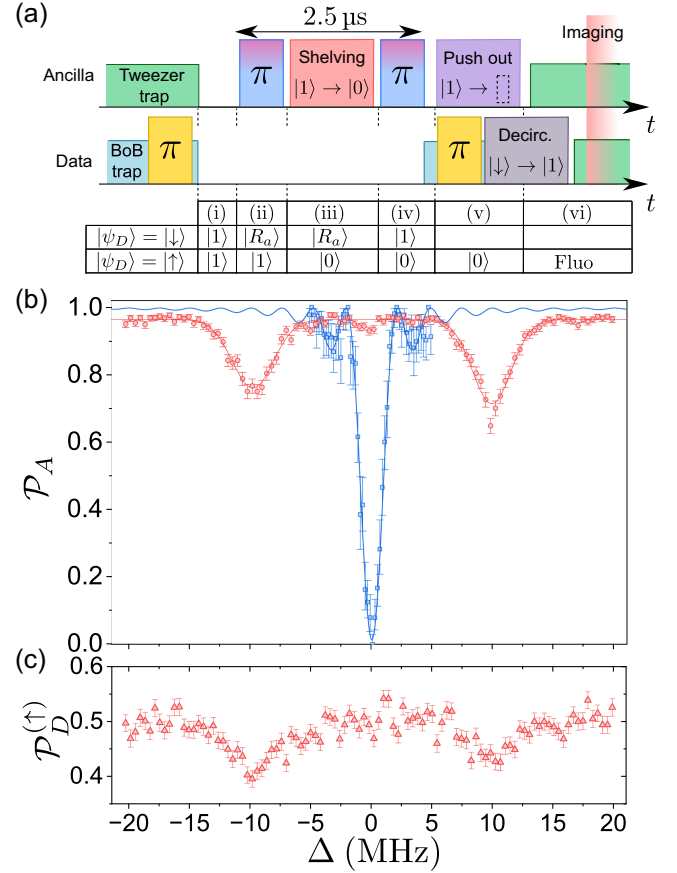


FIG. 2. Ancilla optical excitation spectra. (a) The top (bottom) line is a graphical representation (not to scale) of the sequence of optical (MW) pulses that address the ancilla (data) atom during the ABM sequence. The boxes filled with a blue-red gradient indicate the Rydberg excitation pulses, the green boxes indicate the time during which we trap the ancilla in the optical tweezers. The bottom line displays the time intervals during which the BoB traps are on (cyan boxes), the optional MW π pulses that transfer $|\downarrow\rangle$ to $|\uparrow\rangle$ (yellow boxes), and the decircularization process that brings back the circular state $|\downarrow\rangle$ to $|1\rangle$ (gray box). The table gives the value of $|\psi_A\rangle$ at every step of the sequence in the cases where the MW pulses on the data atom are applied (bottom line) or not (top). The dashed lines indicate the steps of the sequence on the graphical representations. (b) Excitation spectrum of $|R_a\rangle$. We plot as a function of the detuning Δ of the Rydberg excitation lasers from the $|1\rangle \rightarrow |R_a\rangle$ resonance the probabilities $\mathcal{P}_{A|\downarrow}$ (blue) and $\mathcal{P}_{A|\uparrow}$ (red) to recapture the ancilla when the data atom has been prepared in $|\downarrow\rangle$ and $|\uparrow\rangle$, respectively. The solid lines are fits to the data of a single coherent excitation profile (blue) or a sum of two Gaussian peaks (red) (see Appendix B). (c) Probability to recapture the data atom when prepared in $|\uparrow\rangle$. In (b) and (c), the red (blue) points are the result of an average over 400 (50) repetitions, the error bars are statistical (Clopper-Pearson interval).

fluorescence image. The last two steps of the protocol perform a state-selective destructive measurement of the ancilla state $|0\rangle$. Note that we turn off the BoB traps during steps (ii)–(iv) to prevent residual light shifts they could induce on the nearby ancilla.

We use the ABM sequence with a single ancilla to measure the state of the data atom, after its preparation in $|\downarrow\rangle$ or $|\uparrow\rangle$. The effective preparation of these states is checked by a DM of $|\psi_D\rangle$, which we perform after step (iv) of the ABM scheme. The final fluorescence image, which equally detects atoms in $|0\rangle$ or $|1\rangle$, reads the outcome of both the DM and ABM.

In Fig. 2(b), we plot the probabilities to detect the ancilla conditioned on the destructive detection of the data qubit in $|\downarrow\rangle$ and $|\uparrow\rangle$, denoted $\mathcal{P}_{A|\downarrow}$ and $\mathcal{P}_{A|\uparrow}$, respectively, as a function of the detuning Δ of the ancilla optical excitation from the $|1\rangle \rightarrow |R_a\rangle$ transition. The probability $\mathcal{P}_{A|\downarrow}(\Delta)$ (blue squares) displays a coherent excitation profile, which reveals the excitation of the ancilla to $|R_a\rangle$ when $\Delta = 0$, unaltered by the presence of the data atom in $|\downarrow\rangle$. The symmetric splitting of $\mathcal{P}_{A|\uparrow}(\Delta)$ (red circles) around $\Delta = 0$ is a clear signature of the interaction with the data atom in $|\uparrow\rangle$: The excitation to $|R_a\rangle$ at resonance is essentially suppressed. The residual excitation at $\Delta = 0$ is the signature of a false-positive DM of the data atom in $|\uparrow\rangle$, which may stem from a decay to $|0\rangle$ of the low- ℓ Rydberg levels transiently populated during the circular state preparation [32]. The two main peaks correspond to the excitation to $|- \rangle$ or $|+ \rangle$. From their splitting $\delta\nu = 2V/h = 19.6(1)$ MHz, we find an interatomic distance $d = 6.12(1)$ μm close to the 6- μm programmed value [34]. Their widths (4 MHz) are found to be mainly due to the thermal spread of the atomic positions within the traps (see Appendix B). At $\Delta = 0$, the mapping by the ABM sequence of the circular states $|\downarrow\rangle$ and $|\uparrow\rangle$ onto the ancilla states $|0\rangle$ and $|1\rangle$ is conspicuous: The ancilla is detected in $|0\rangle$ only if the data atom has been prepared in $|\uparrow\rangle$. From the data of Fig. 2(b), we estimate the false-negative probability $1 - \mathcal{P}_{A|\uparrow}(\Delta = 0)$ to be about 6.7% and false-positive probability $\mathcal{P}_{A|\downarrow}(\Delta = 0) \approx 1.3\%$. This corresponds to a QND readout fidelity $\mathcal{F} = (\mathcal{P}_{A|\uparrow} + 1 - \mathcal{P}_{A|\downarrow})/2$ of 96%, significantly improving over previous demonstrations with dual-species experiments [25].

In the case where we prepare the data atom in $|\uparrow\rangle$, we plot in Fig. 2(c) the probability $\mathcal{P}_D^{(\uparrow)}$ to recapture the data atom after the final DM of $|\uparrow\rangle$. The probabilities are the same whether $\Delta = 0$ (full ABM) or $|\Delta| > 15$ MHz (no ABM), showing that the state of the data atom is unaltered by the ABM. The reduction of $\mathcal{P}_D^{(\uparrow)}$ when $\Delta \sim \pm V/h$ corresponds to the excitation of the $|\pm\rangle$ states. A fraction of the data state then overlaps with $|R_c\rangle$ and is not detected by the DM.

We verify that the ABM fulfils the requirements of a genuine QND measurement by using it to record Rabi oscillations of the data atom. We initially prepare an ensemble with two ancillae in the $|1, \downarrow, 1\rangle_{ADA}$ state. We apply an MW pulse of variable duration τ [see Fig. 3(a)], resonant with the $|\downarrow\rangle \rightarrow |\uparrow\rangle$ transition, and measure the state $|\uparrow\rangle$ with the DM and the ABM schemes. We plot in

Fig. 3(b), as a function of τ , in blue $\mathcal{P}_D^{(\uparrow)}$ and in red the probability to recapture both ancillae, denoted $\mathcal{P}_{AA}$. Rabi oscillations are apparent in both signals, showing that the ABM reproduces the evolution of $|\psi_D\rangle$ found by the DM. Remarkably, the contrast of the ABM scheme is notably higher than that of the DM scheme. The ABM is not affected by the 66% efficiency of the transfer from $|\downarrow\rangle$ to $|1\rangle$, which corresponds to the detection efficiency of the DM (see Appendix C). The ABM contrast of 67.3(8)% is essentially limited by the 72.3(4)% preparation efficiency of $|\downarrow\rangle$ rather than by the ABM detection efficiency. Note that the preparation efficiency could be significantly improved by operating in a cryogenic setup and by employing optimal control, as already demonstrated in Ref. [35].

To assess the ABM fidelities without contamination of the DM level crosstalks which affect previous estimations from Fig. 2 (see Appendix C), we measure the joint probabilities to detect none, part of, or all of the three atoms in the fluorescence imaging that follows the Rabi oscillations, for two sequences in which we perform a DM of $|\uparrow\rangle$ [as in Fig. 3(b)] or $|\downarrow\rangle$. We define $\eta_\uparrow$ ($\eta_\downarrow$) as the probability for the ancilla to stay in $|1\rangle$ (to be transferred to $|R_a\rangle$), at step (ii) of the ABM sequence, when the data qubit is in $|\uparrow\rangle$ ($|\downarrow\rangle$). From an experimental-imperfection model fitted to the measured Rabi oscillations (see Appendix C), we get $\eta_\uparrow = 96.6(2)\%$ and $\eta_\downarrow = 95.0(3)\%$. The efficiency $\eta_\uparrow$ is limited by ancilla losses due to collisions with the background gas or to finite recapture efficiency. The main limitations to $\eta_\downarrow$ are the finite excitation efficiency of $|R_a\rangle$ for an isolated atom and the lifetime of $|R_a\rangle$. These technical limitations may be reduced in an improved setup with, e.g., a cryogenic environment to increase the lifetimes of both the atoms in their traps and of the Rydberg levels.

Finally, we check that the ABM possesses the projective character of a QND measurement. After an MW $\pi/2$ pulse that prepares the data atom in $(|\downarrow\rangle + |\uparrow\rangle)/\sqrt{2}$, we perform the ABM sequence with a single ancilla, apply a second variable-duration MW pulse that drives Rabi oscillations [see Fig. 3(c)], and finish with the DM of the data atom. Discarding the ABM information, we get $\mathcal{P}_D^{(\downarrow)}(\tau)$ and plot it in Fig. 3(d). In spite of the second MW pulse, we find no oscillations. However, if we select only the experimental realizations where the ABM has measured the data atom in $|\uparrow\rangle$, we recover fully contrasted oscillations of $\mathcal{P}_D^{(\downarrow)}(\tau)$, starting from $\mathcal{P}_D^{(\downarrow)}(0) = 0.01(1)$, revealing the maximal correlation between the data and the ancilla states. The ABM has projected the data atom onto $|\uparrow\rangle$, as expected for a genuine QND measurement. Remarkably, this also shows that the ABM can be used to measure $|\psi_D\rangle$ at any time in the sequence and thus to operate midcircuit measurements. The duration of the QND measurement is ultimately limited by the 2.5- μs duration of steps (ii)–(iv), used to entangle the data atom and the ancilla atom while leaving the latter in its ground state. The slow readout of the measurement

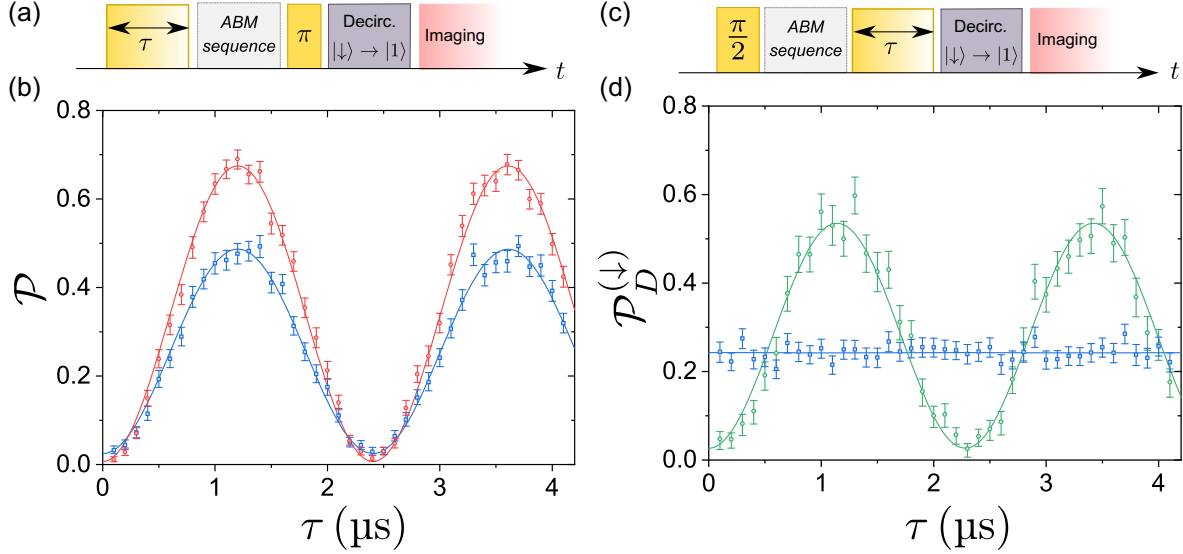


FIG. 3. Assessing the QND character of the ABM with Rabi oscillations. (a) Pulse sequence used to record Rabi oscillations on the $|\downarrow\rangle \rightarrow |\uparrow\rangle$ transition. After an MW pulse of duration τ , the population in $|\uparrow\rangle$ is measured with the ABM and the DM schemes. The color code is that of Fig. 2(a). (b) Probabilities $\mathcal{P}_{AA}$ and $\mathcal{P}_D^{(\uparrow)}$ to recapture the ancillae (red circles) and the data atom (blue squares), respectively, measured with the sequence sketched in (a). The solid lines are sinusoidal fits to the data with a shared oscillation frequency and fitted contrasts of 0.668(7) and 0.461(4) for the ABM and DM, respectively. (c) Pulse sequence used to record Rabi oscillations with the DM, set to measure $|\downarrow\rangle$. The oscillations are driven after a $\pi/2$ pulse that prepares the data atom in $(|\uparrow\rangle + |\downarrow\rangle)/\sqrt{2}$ and an ABM of the data state. The color code is that of Fig. 2(a). (d) Probability $\mathcal{P}_D^{(\downarrow)}$ to recapture the data atom averaged over all realizations of the experiment (blue squares), i.e., unconditioned on the ABM outcome, or conditioned on the recapture of the ancilla, i.e., on the nondestructive measurement of $|\uparrow\rangle$ (green circles). The solid lines are sinusoidal fits to the data with shared oscillation frequencies and fitted contrasts of 0.51(1) and $0.5(50) \times 10^{-4}$. In (b) and (d), the data points are the result of an average over 200 realizations of the experiment, the error bars are statistical.

outcome, corresponding to the 25-ms-long fluorescence imaging, is performed at the end of the experimental sequence. It does not limit the characteristic timescales of phenomena that could be observed through the ABM. Note that the fluorescence measurement could be performed in a microsecond timescale [16], too, and be made nondestructive [36]. This would be particularly important for error correction, which requires the measurement to be read midcircuit.

We now use the ancilla to optically apply phase and spin flips to the data atom. An optical 2π pulse on the $|1\rangle \rightarrow |R_a\rangle$ ancilla transition maps the states $|\downarrow, 1\rangle_{DA}$ and $|\uparrow, 1\rangle_{DA}$ to $-|\downarrow, 1\rangle_{DA}$ and $|\uparrow, 1\rangle_{DA}$, respectively. When $|\psi_D\rangle = |\uparrow\rangle$, the Rydberg blockade of the ancilla excitation [Fig. 1(b)] inhibits the π phase shift experienced by $|\downarrow, 1\rangle_{DA}$, which we probe with an MW Ramsey interferometer [see Fig. 4(a)]. In Fig. 4(b), we plot $\mathcal{P}_D^{(\downarrow)}$ measured after two MW $\pi/2$ pulses on the $|\downarrow\rangle \rightarrow |\uparrow\rangle$ transition separated by 3 μ s, as a function of the detuning δ of the driving MW field from resonance. The Ramsey fringes are measured either with (green) or without (blue) the optical 2π pulse on the ancilla. The expected π phase shift is conspicuous. A sine fit to the data reveals a $1.01(2)\pi$ phase shift with a contrast reduction by a factor of only 0.93(4). This demonstrates our ability to optically imprint a π phase shift on the circular-state-encoded data qubit, naturally impervious to optical fields.

Moreover, when $\delta = 0$, this protocol realizes an ancilla-controlled spin flip of the data qubit, demonstrated here with the atoms initially in $|\downarrow\rangle$, with a fidelity corresponding to the fringe contrast reduction, i.e., 93%.

We also demonstrate an optically-controlled spatially-resolved state transfer of the data atom by using an ancilla excited to $|R_a\rangle$, which blocks the MW transition from $|\downarrow\rangle$ to $|\uparrow\rangle$ [Fig. 1(b)]. We prepare an array containing four data atoms, initially excited to $|\downarrow\rangle$, each coupled to two ancillae, as drawn in the inset of Fig. 4(d). One of the ancillae is used for the control of the data atom state, and both of them for the measurement of this state. We block the transition from $|\downarrow\rangle$ to $|\uparrow\rangle$, induced by an MW π pulse [see Fig. 4(c) for the pulse sequence], in the two ensembles on the right by exciting the control ancillae to $|R_a\rangle$ before the MW pulse. To do so, we turn off the optical tweezers during the excitation to $|R_a\rangle$ for these two ancillae only (see Appendix A). The six other ancillae remain in $|1\rangle$ because of the trap-induced light shifts. After the MW pulse, a second optical π pulse brings the two control ancillae back to $|1\rangle$. We eventually turn back on all the tweezer traps. This leaves the two right (left) data atoms in $|\downarrow\rangle$ ($|\uparrow\rangle$) and all ancillae in $|1\rangle$. Remarkably, the ancillae, including the control ones, can be used again to perform a QND measurement of the state of the data atom. We plot in Fig. 4(d) $\mathcal{P}_{AA}$ for the four ensembles of the array, as a

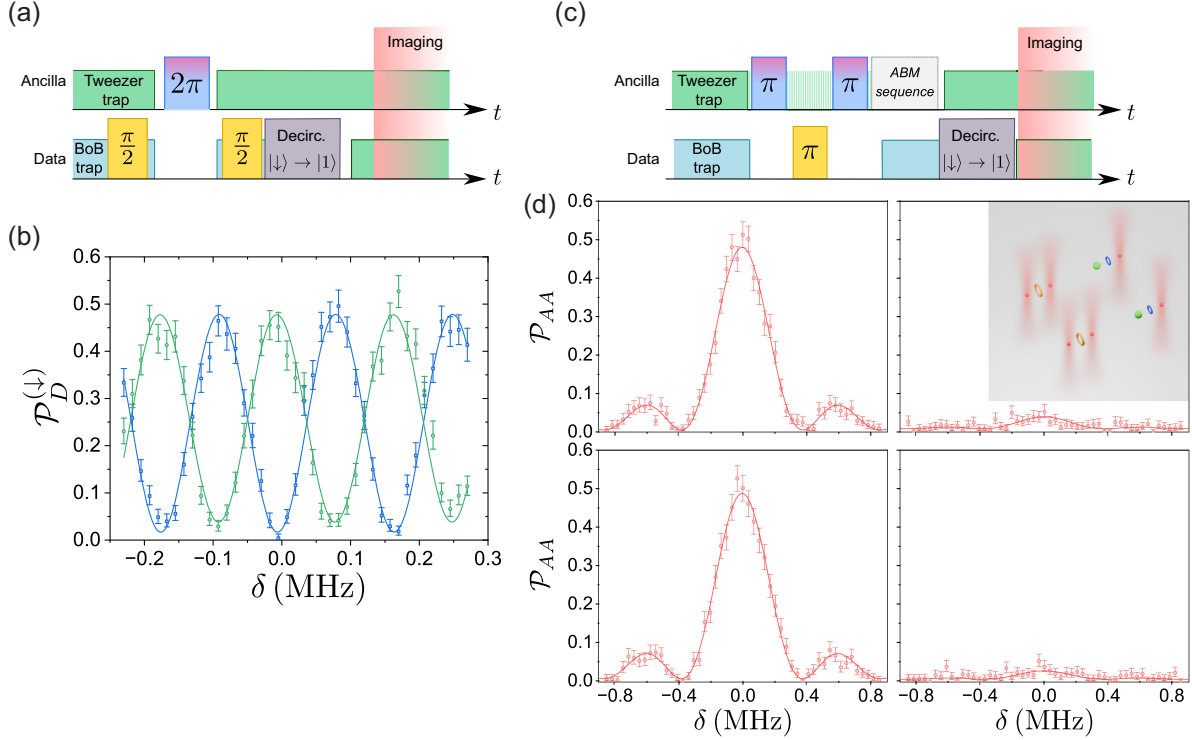


FIG. 4. Local manipulation of the data atom with the ancilla. (a) Pulse sequence used to measure a π phase shift on the data atom using a 2π optical pulse on the ancilla. The two MW $\pi/2$ pulses define a Ramsey interferometer. The color code is that of Fig. 2(a). (b) Probability to measure the data atom in $|\downarrow\rangle$ with the destructive scheme after the Ramsey sequence of (a), as a function of the frequency of the MW field. The optical 2π pulse on the ancilla atom is either applied (green circles) or absent (blue squares). The points are the result of an average over 100 repetitions, the error bars are statistical. The solid lines are fits of a sine law to the data with contrasts 0.46(1) (green) and 0.44(1) (blue) and a $1.01(2)\pi$ phase difference. (c) Pulse sequence used to locally manipulate the state of the data atom. We apply an MW π pulse on the data atom in between two optical π pulses on the ancillae, during which the tweezer traps are turned off (striped green area) for the control ancillae. (d) Inset: structure of the array of four ensembles of data and ancilla atoms. The two ancillae, the tweezer traps of which are turned off during the optical π pulses of (c), are excited to $|R_a\rangle$ (green sphere) while the others stay in $|1\rangle$ (red sphere). The two corresponding data atoms on the right stay in $|\downarrow\rangle$ (blue torus) while the two left ones are transferred to $|\uparrow\rangle$ (yellow torus). Probability (red circles) to measure the data atom in $|\downarrow\rangle$ with the ABM, plotted as a function of the frequency of the MW field. Each panel corresponds to the ensemble with the corresponding position in the inset. The points are the result of an average over 400 repetitions, the error bars are statistical. The solid red lines are fits of a spectral line profile to the data (see Appendix B), with amplitudes, from left to right and top to bottom, of 0.48(1), 0.03(1), 0.485(7), and 0.022(5).

function of the MW frequency detuning δ , exhibiting the excitation spectrum of the $|\downarrow\rangle \rightarrow |\uparrow\rangle$ transition. While the left atoms end up in $|\uparrow\rangle$ when $\delta = 0$, the MW excitation of the right atoms is strongly suppressed. The residual excitation results from the finite efficiency of the initial excitation of the ancilla to $|R_a\rangle$ (see Appendix C). We quantify the site selectivity of the state transfer by calculating $\mathcal{S} = 1 - (A_r/A_l)$, where $A_{r,l}$ is the average value of the amplitudes of the excitation spectra for the right (r) or left (l) atoms, and find $\mathcal{S} = 95(2)\%$.

We have developed a hybrid dual-Rydberg platform that enables high-fidelity quantum nondemolition measurements and spatially-resolved optical manipulations of individual circular Rydberg atoms. The entanglement time is in the microsecond range, well matched to the timescales of the circular-Rydberg-atom platforms [28,37]. It could be reduced by replacing steps (ii)–(v) of the ABM sequence with a 2π pulse on the $|1\rangle \rightarrow |R_a\rangle$ transition, which

effectively operates a controlled-Z gate between the data atom and the ancilla, within a Ramsey interferometer on the ground-state hyperfine transition [28]. The interferometer would read the phase shift and, thus, the state of the data atom. This would in addition reduce the imperfections due to the finite lifetime of the low- ℓ $|R_a\rangle$ state. This would also provide a measurement of the entanglement fidelity, not performed in this work.

In our experiment, we use up to two ancillae per data atom. Using an assembly of ancilla atoms could increase the measurement fidelity [23]. It would also enable the detection of multiple circular Rydberg levels or the measurement of the same atom at different times. The ancillae could also be used to mediate rapidly tunable interactions between data atoms [38], even at a distance, hence increasing the system connectivity [39]. Finally, our method is not restricted to the two circular Rydberg states employed in this work (see Appendix B). It can be readily

extended to quantum simulators that involve other non-circular high-angular-momentum Rydberg levels [40], using the wide Stark tunability of Rydberg levels to reach a Förster resonance, or circular Rydberg states of other species than alkali atoms, e.g., strontium [41,42].

This hybrid dual-Rydberg platform is very promising for quantum simulation. It allows for the initialization of the simulator in an arbitrary product state. The QND nature of the detection allows one to measure correlations between arbitrary single-qubit observables at different times, including out-of-time-order correlations [43–46]. This is a crucial feature to take full benefit of the long simulation times enabled by circular Rydberg atoms [37]. It opens the route to the study of long-term coherent evolutions, slow thermalization dynamics [47,48], and to the observation of measurement-induced phase transitions [49–51].

This work has received funding by the France 2030 programs of the French National Research Agency (Grant No. ANR-22-PETQ-0004, project QuBitAF), under Horizon Europe programme HORIZON-CL4-2022-QUANTUM-02-SGA via Project No. 101113690 (PASQuanS2.1), by the European Union (ERC Advanced Grant No. 786919, project TRENSCRYBE). It has been supported by the Île-de-France region in the framework of DIM QuanTiP (project LT-CRAQS) and by the Quantum Information Center Sorbonne as part of the program *Investissements d'excellence*—IDEX of the Alliance Sorbonne Université.

DATA AVAILABILITY

The data that support the findings of this article are openly available [52].

APPENDIX A: EXPERIMENTAL PROCEDURES

In Appendix A, we describe the experimental setup, the arrays of trapping sites that we use in the experiment, and the main steps of an experimental sequence.

1. Experimental setup

The experimental setup is described in detail in Ref. [32] and is sketched in Fig. 5. The atoms are trapped inside a hollow sapphire cube, which is held under ultrahigh vacuum. The cube supports electrodes to dynamically control the dc or radio-frequency electric fields and two tightly focusing lenses, L_1 and L_2 , used for the generation of the trapping sites (see Appendix A2). The Rydberg-excitation laser beams and an additional laser beam, used to optically pump and push out ground-state atoms, enter the cube along the x axis (see Fig. 5 for axes definition). All trapping laser beams enter along the y axis and are focused by lens L_1 onto the atoms. This same lens collects the fluorescence photons emitted by the trapped atoms, and

directs them toward an electron-multiplying CCD camera (EMCCD) camera.

2. Generation of the trapping sites

We prepare three different arrays of trapping sites by holography, using three liquid-crystal spatial light modulators (SLM) [Hamamatsu LCOS X10468 (SLM1 and 3) and X15213 (SLM2)]. An SLM imprints a digitally-controllable phase mask on the impinging collimated Gaussian beam. We program a phase mask so that its diffraction pattern, obtained in the focal plane of the tightly focusing lens L_1 (focal length $f = 16.32$ mm, numerical aperture $NA = 0.27$), corresponds to the desired array of trapping sites. We use the SLMs, in addition, to correct the optical aberrations of the setup. The method we use to generate the phase masks is described in Refs. [32,53,54].

The three SLMs are arranged as sketched in Fig. 5. The geometry of the trapping sites is sketched in Fig. 6. All the trapping beams are generated from a 2-W laser of 821-nm wavelength (Toptica TA Pro). The three beams are overlapped with polarizing beam splitter cubes and can be turned on or off independently using acousto-optic modulators (AOM) not shown in Fig. 5.

- (i) SLM1 prepares an array of $13 \times 5 = 65$ Gaussian optical tweezers with a waist of $1.2 \mu\text{m}$, with an intersite spacing of 6 and $12.5 \mu\text{m}$ along the x and z axes, respectively. With a power of 2.6 mW per trapping site, the transverse and longitudinal trap oscillation frequencies are 80 and 17 kHz, respectively.
- (ii) SLM2 prepares an array of four optical BoBs, separated by 30 and $25 \mu\text{m}$ along the x and z axes, respectively. The power per BoB is 22 mW, corresponding to transverse and longitudinal trap oscillation frequencies of 16 and 5 kHz, respectively.
- (iii) SLM3 is used to generate either a single or an array of Gaussian optical tweezers. The shape of the optical tweezers is identical to that of those prepared by SLM1. This array of optical tweezers is meant to trap the ancilla atoms, while we use the single optical tweezer for the preparation of defect-free arrays of atoms. A telescope conjugates the plane of SLM3 with that of an acousto-optic deflector (AOD), so that a dynamically controllable phase grating generated by the AOD is superimposed with the phase mask generated by SLM3. We use the AOD during an experimental sequence to move either a single (see Appendix A3) or an array (see Appendix A5) of Gaussian optical tweezers.

3. Preparation of the ensembles of trapped atoms

We initially load atoms into the array of optical tweezers prepared by SLM1, in the collisional blockade regime. We employ a 3D magneto-optical trap (MOT), initially loaded by a low-velocity atomic beam, originating from a 2D

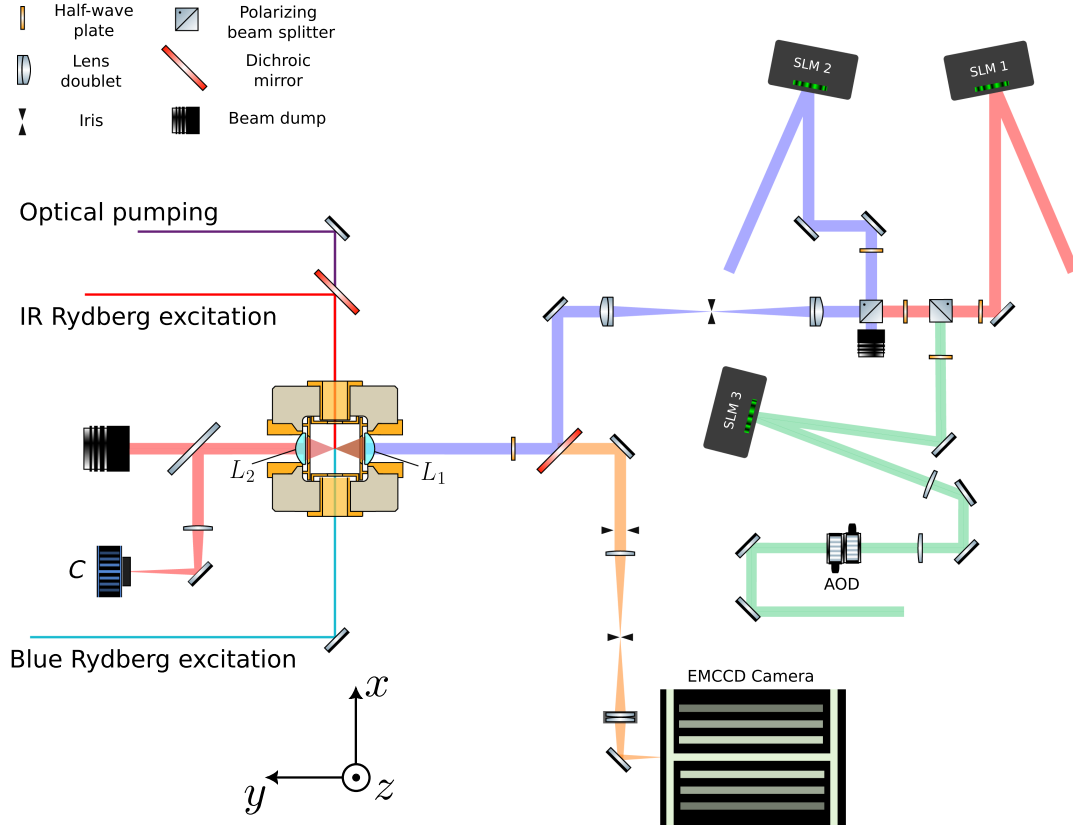


FIG. 5. Sketch of the experimental setup. A sapphire cube (beige, not to scale) holds two short-focal-length lenses (cyan) and electrodes (orange) that control the electric field. The blue and infrared Rydberg excitation lasers (thin light blue and red lines) and the ground-state manipulation beam (thin violet line) enter the cube along the x axis. The three trapping light beams (red, violet, and green) prepared by the three SLMs and overlapped with polarizing beam splitter cubes enter the cube along the y axis and through lens L_1 . The AOD is used to dynamically move the set of tweezers prepared by SLM3. A second lens L_2 collects all trapping light, eventually focused on a monitoring camera (C). Fluorescence photons (orange) are collected by L_1 and focused into the EMCCD camera. Several irises spatially filter the light to reduce stray light. The SLM planes are imaged onto lens L_1 with a $4f$ telescope with $f = 400$ mm. An iris placed at the intermediate focal point of this telescope filters out nondiffracted light.

MOT, that enters the vacuum chamber along the z axis. A subsequent optical-molasses stage cools the trapped atoms to 13 μ K. A 25-ms-long fluorescence image, using the optical-molasses beams, detects the trapped atoms. We measure an average loading probability per site of 62%. We then use the single Gaussian optical tweezer generated by SLM3 to rearrange the atoms into the target array [3], indicated in Fig. 6(a) by red circles. This array consists in four ensembles of either two or three atoms, depending on the number of ancillae per ensemble. We use a second fluorescence image to probe the effective preparation of a defect-free target array. All the data presented in this work are postselected to keep only the experimental realizations for which the target array had no defects, and no additional atoms were left in the other trapping sites of SLM1.

After the second fluorescence image, we adiabatically lower the power of the trapping sites of SLM1 to 0.16 mW, cooling the atoms down to 4.5 μ K. We increase the

magnetic field to 14 G and optically pump all atoms to $|1\rangle = |5S_{1/2}, F=2, m_F=2\rangle$ using a σ^+ polarized laser beam on resonance with the $|5S_{1/2}, F=2\rangle \rightarrow |5P_{3/2}, F'=2\rangle$ transition, in conjunction with a repumper laser beam on resonance with the $|5S_{1/2}, F=1\rangle \rightarrow |5P_{3/2}, F'=2\rangle$ transition. Simultaneously, we change the phase mask applied on SLM3, so that it now prepares the array of trapping sites shown in Fig. 6(c). We adiabatically turn on this array of trapping sites for the ancilla atoms up to 1.0 mW per site within 0.5 ms. At this stage, the atoms that will be used as ancillae are trapped in the superposition of two Gaussian optical tweezers, created by SLM1 and SLM3. Because of the large frequency difference (≈ 200 MHz) between the two laser beams and because of the significant intensity difference, however, the trapping potential is essentially controlled by the optical tweezers created by SLM3. The data atoms are trapped in the low-power optical tweezers from SLM1 only.

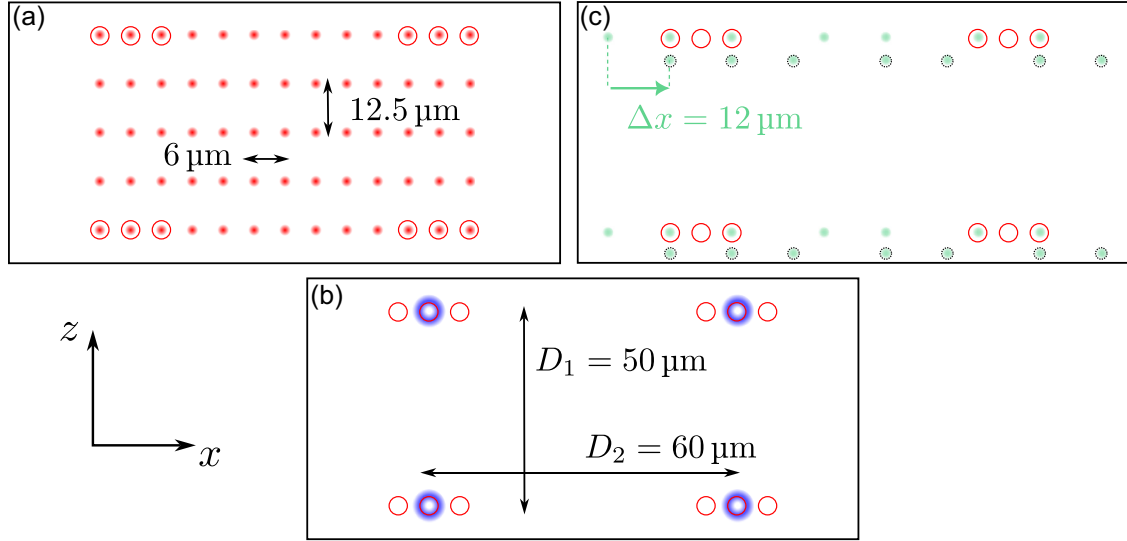


FIG. 6. Geometries of the three arrays of trapping sites. (a) Arrays of Gaussian optical tweezers (red gradient filled circles) generated by SLM1. The 65 sites are arranged along a regular array with an intersite spacing of 6 and 12.5 μm along the x and z axes, respectively. The axis definition is identical to that of Fig. 5. In all panels, the red open circles indicate the 12 sites of the target array, deterministically prepared with the rearranger beam. For the experiments with one ancilla per ensemble, the right-hand sites of the ensembles are left empty. (b) Array of four optical bottle beams (violet rings), prepared by SLM2, that are used to trap the circular Rydberg atoms. (c) Array of Gaussian optical tweezers (green gradient filled circles) prepared by SLM3, used to trap the ancilla atoms. The circles with dark dashed contour lines indicate the position of the trapping beams when the AOD shifts the whole array by $\Delta x = 12 \mu\text{m}$. For clarity, they are shifted down along the z axis with respect to their actual positions. In this situation, there is no remaining light in two out of the eight original sites.

4. Experimental sequence

After the preparation stage described in the previous section, we run an experimental sequence, of which we sketch a typical example in Fig. 7. The latter corresponds to the sequence of Fig. 4(c) and the data of Fig. 4(d). At $t = 0$, we turn off the optical tweezers from SLM1, leaving the data atoms untrapped. The ancillae are still trapped in the Gaussian tweezers of SLM3. The data atoms are then excited to $|52D_{5/2}, m_J = 5/2\rangle$ using a 1- μs -long two-photon optical π pulse with σ^+ polarized laser beams of wavelengths 420 and 1015 nm, detuned by 500 MHz from the intermediate $|6P_{3/2}, F' = 3\rangle$ level. The ancillae are impervious to this excitation because of the 8.6-MHz trap-induced light shifts. We turn on the BoBs prepared by SLM2 right after the optical excitation.

An MW field at 64.7566 GHz transfers the data atoms to $|52F, m = 2\rangle$. We adiabatically increase the electric field to 2.16 V cm^{-1} while shining a σ^+ polarized radio-frequency field at 225 MHz to transfer these atoms to the circular state $|\downarrow\rangle = |52C\rangle$. Then, we set the electric field to 1.59 V cm^{-1} to reach the required Förster resonance (see Appendix B) before running one of the sequences of MW and optical pulses described in the main text.

At the end of the sequence, all atoms in any ground state (including $|0\rangle$ and $|1\rangle$) are recaptured in the array of trapping sites prepared by SLM1. The two other arrays are turned off. We use a third fluorescence image, identical

to the previous ones, to measure the probability to recapture the ancilla and data atoms in their respective optical tweezers. To reduce the probability that an atom in a low- ℓ Rydberg state at the end of the sequence decays back to its ground state and is detected by the final fluorescence image, we increase the electric field to 8.2 V cm^{-1} after the last optical pulse that transfers the $|52D\rangle$ state back to $|1\rangle$. This strongly reduces the decay rates on optical transitions from Rydberg levels to the ground states, so that Rydberg atoms remain as long as possible in antitrapped Rydberg levels and are expelled from the Gaussian optical tweezers well before the fluorescence imaging.

5. Manipulation of the ancillae

The optical excitation of the ancillae from $|1\rangle$ to $|R_a\rangle = |45S_{1/2}, m_J = 1/2\rangle$ is similar to that of the data atoms. It employs the same 420-nm-wavelength light, together with a σ^- polarized 1017-nm-wavelength light, which is overlapped with the 1015-nm-wavelength beam used for the data atoms. The optical π pulse that prepares $|R_a\rangle$ lasts 0.4 μs .

We turn off the BoBs during the optical excitation of the ancillae, from $|1\rangle$ to $|R_a\rangle$ or back, to avoid residual light shifts. If all of the ancillae are to be excited, we turn off the laser beam impinging on SLM3 using an AOM (not shown in Fig. 5). For the data of Fig. 4(d), however, only two out of eight ancillae should be excited to $|R_a\rangle$. To do this, using

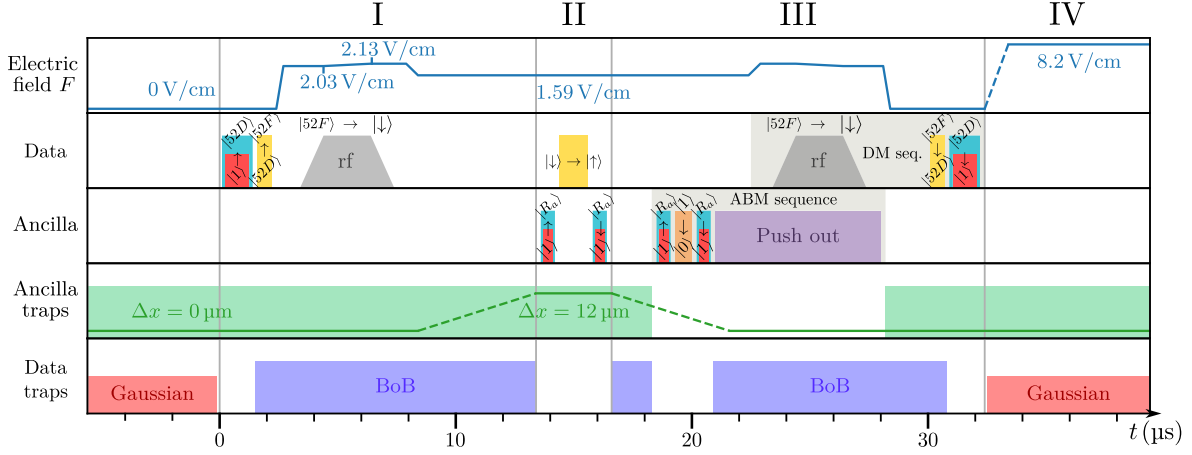


FIG. 7. Typical experimental sequence run after the preparation of a defect-free array of individual atom, which ends at $t = 0$. The sequence is schematically represented as a function of time. Its four main parts at $t > 0$ correspond to (I) the preparation of the data atom in $|\downarrow\rangle$, (II) local manipulation of the data atom that corresponds to the data of Fig. 4(d), (III) the ABM and DM sequences used for the nondestructive and destructive measurement of the data atom, and (IV) the final readout of the state with a fluorescence image (not shown) once the atoms have been recaptured in the optical tweezers of SLM1. Right after the last Rydberg excitation pulse, we increase the electric field to reduce the decay of residual Rydberg excitations to the ground state. First line: amplitude of the electric field. Second line: optical (blue and red), MW (yellow), and rf (gray) pulses applied to the data atom. The trapeze form of the radio-frequency pulse corresponds to an adiabatic increase and reduction of the rf power. The transitions addressed by each pulse are superimposed with the color boxes. The decaricularization process used for the DM is indicated by a light gray background. Third line: same as the second line, but for the ancilla atoms. The orange and violet boxes correspond to the shelving and push-out stages, respectively, of the ABM sequence indicated by a light gray background. Fourth line: the green color box indicates when the trapping light from SLM3 is applied onto the ancilla atoms. The green line is a plot of the programmed shift of the whole array imposed by the AOD. The dashed line corresponds to the finite response time of the AOD during which the trapping array is changing. Fifth line: the red and blue boxes indicate when the trapping light from SLM1 (Gaussian) or SLM2 (BoB), respectively, is shone onto the atoms.

the AOD, we translate the array of trapping sites generated by SLM3 along the x axis by $\Delta_x = 12 \mu\text{m}$. Using the geometry shown in Fig. 6, this translation effectively leaves only two ancillae untrapped, while the six other are still trapped after the displacement. The latter lasts $5 \mu\text{s}$, limited by the response time of the AOD. We find that we do not need to turn off the trapping light during the translation of the array to avoid atom losses.

The selective detection of state $|1\rangle$ after step (ii) of the ABM sequence requires one to shelve state $|1\rangle$ into $|0\rangle$ [step (iii) of the ABM sequence]. We perform the shelving into $|0\rangle$ with a $0.7\text{-}\mu\text{s}$ -long pulse of optical pumping on the π -polarized $|1\rangle \rightarrow |5P_{3/2}, F' = 2, m_F = 2\rangle$ transition. The corresponding laser beam enters the vacuum chamber with a 45° angle from the y and z axes (not shown in Fig. 5) and has an intensity 170 times higher than the saturation intensity of the transition. The atoms excited to $|R_a\rangle$ at step (ii) are deexcited to $|1\rangle$ at step (iv) of the ABM sequence and expelled from the traps at step (v), without affecting those in $|0\rangle$. We use a σ^+ polarized push-out laser beam at resonance with the $|1\rangle \rightarrow |5P_{3/2}, F' = 3, m_F = 3\rangle$. Its intensity is 0.66 times the saturation intensity of the transition and we shine it during $7 \mu\text{s}$. It shares the same optical path as the laser used for the initial optical pumping into $|1\rangle$. During both the shelving and the push-out phases, we turn off the trapping laser beams of the ancillae.

APPENDIX B: STARK-TUNED FÖRSTER RESONANCE

In Appendix B, we discuss the choice of levels to achieve a Förster resonance.

1. Choice of the Rydberg states

The QND measurement and the local manipulation of the data atoms rely on the dipole blockade of optical or MW transitions in the presence of interactions between the data atom in a circular Rydberg level $|\uparrow\rangle = |nC\rangle$ and an ancilla in a laser-accessible low-angular-momentum (ℓ) Rydberg level $|R_a\rangle$. Strong first-order interactions between $|\uparrow\rangle$ and $|R_a\rangle$ require a Förster resonance. However, there is *a priori* no zero-field resonance for Rb between a pair of high- ℓ Rydberg levels $|R_c\rangle, |nC\rangle$ and a couple of low- ℓ Rydberg states $|R_a\rangle, |R'_a\rangle$, with $|R_a\rangle$ being laser accessible from the ground state and $|R_a\rangle \rightarrow |R'_a\rangle$ and $|R_c\rangle \rightarrow |nC\rangle$ being electric-dipole-allowed transitions. Therefore, we rather rely on a Stark-tuned Förster resonance [55,56]: We seek pair states $|R_a, R'_a\rangle$ and $|nC, R_c\rangle$ at resonance in an experimentally convenient electric field of a few V cm^{-1} .

In order to maximize the strength of the interaction, the electric dipole matrix elements of the $|R_a\rangle \rightarrow |R'_a\rangle$ and $|R_c\rangle \rightarrow |nC\rangle$ transitions should be as large as possible, leading to the choice of $|R_c\rangle = |(n \pm 1)C\rangle$. In addition,

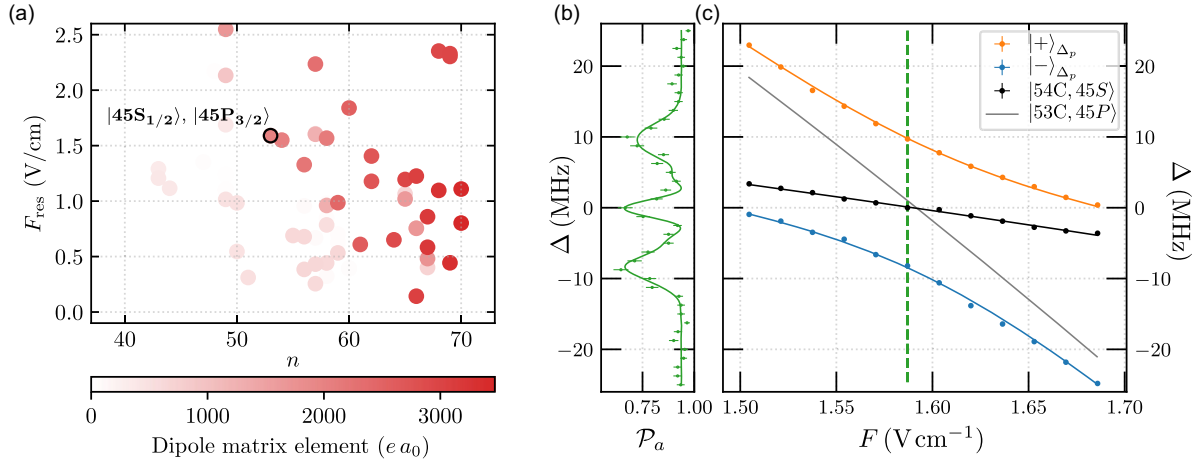


FIG. 8. Choice and spectral evidence of a Förster resonance. (a) Map of the low-angular-momentum Rydberg pair states $|R_a, R'_a\rangle$ that have a Stark-tuned Förster resonance with a circular-Rydberg pair state $|nC, (n+1)C\rangle$. Each point corresponds to a resonance, plotted as a function of n and the resonance electric field F_{res} . The color scale gives the dipole matrix element of the $|R_a\rangle \rightarrow |R'_a\rangle$ transition in units of (ea_0) . The pair state used for our experiments is highlighted by a dark circle. (b) Probability $\mathcal{P}_A$ to recapture an ancilla atom excited to $|45S_{1/2}, m_J = 1/2\rangle$ in the vicinity of a data atom prepared in $|54C\rangle$ in an electric field of $F = 1.591 \text{ V cm}^{-1}$. The experimental sequence is similar to that used for the data of Fig. 2 of the main text. The green solid line is a fit to the data of a sum of two Gaussian peaks and of a coherent excitation profile [Eq. (B4)]. (c) Fitted resonant frequencies of the three peaks observed in the excitation spectrum of (b) plotted as a function of F . Orange, blue, and black data points correspond to the highest, lowest, and central measured frequencies, respectively. They correspond to the frequencies of the transitions from $|54C, 45S\rangle$ to $|+\rangle_{\Delta_p}$ (orange points), $|-\rangle_{\Delta_p}$ (blue points) with interacting atoms, or to $|54C, 45S\rangle$ (black points) for noninteracting atoms. The black line is a fit of Eq. (B5) to the black data points, the orange and blue lines are fits of Eq. (B6) to the orange and blue data points. The light gray line indicates the energy difference between the noninteracting pair states $|54C, 45S\rangle$ and $|53C, 45P\rangle$. An anticrossing, signature of the Förster resonance between the ancilla and data atoms, is apparent. The green dashed line indicates the electric field used in the experiments of the main text.

we restrict ourselves to the situation where the electric field is parallel to the interatomic axis. The conservation of angular momentum then imposes the polarization of the $|R_a\rangle \rightarrow |R'_a\rangle$ to match that of the $|n\pm 1C\rangle \rightarrow |nC\rangle$ transition, i.e., to be σ^+ . In Fig. 8(a), we represent as a function of n and of the resonance electric field F_{res} all Förster resonances we have found between a circular-Rydberg-level pair state $|nC, (n+1)C\rangle$, $43 \leq n \leq 70$, and a low- ℓ Rydberg pair state $|R_a, R'_a\rangle$, where $|R_a\rangle$ and $|R'_a\rangle$ are S , P , or D states, and for which $|R_a\rangle \rightarrow |R'_a\rangle$ is σ^+ polarized. The color code indicates the strength of the dipole matrix element of the $|R_a\rangle \rightarrow |R'_a\rangle$ transition. When different pair states differ only through the quantum numbers m or m_J , we plot only the pair states with the maximum m value, i.e., with the largest dipole.

In our setup, because of MW and rf bandwidth limitations, we cannot efficiently manipulate circular states with $n \gtrsim 55$. We also prefer to work at low field values to reduce the decoherence induced by electric-field noise, though high enough to preserve the stability of circular states. This led us to choose the states $|R_c\rangle = |53C\rangle$, $|R_a\rangle = |45S_{1/2}, m_J = +1/2\rangle \equiv |45S\rangle$, and $|R'_a\rangle = |45P_{3/2}, m_J = +3/2\rangle \equiv |45P\rangle$. The pair states $|53C, 45P\rangle$ and $|54C, 45S\rangle$ are on resonance in a field of $F_{\text{res}} = 1.598 \text{ V/cm}$. The resonant spin-exchange-interaction rate between these levels is $\Omega = 2V/h = 21 \text{ MHz}$, when the data and ancilla atoms are $6.0 \mu\text{m}$ away.

2. Förster resonance spectroscopy

We probe the Förster resonance between the ancilla and data atoms by measuring, for different values of the applied electric field, the laser excitation spectrum of the $|1\rangle \rightarrow |R_a\rangle = |45S_{1/2}, m_J = +1/2\rangle$ transition, when the data atom is prepared in $|\uparrow\rangle = |54C\rangle$. The experimental sequence is similar to that used for the data shown in Fig. 2(b) of the main text, but the Rydberg-excitation optical pulse lasts here $0.3 \mu\text{s}$. The spectrum shown in Fig. 8(b) is measured close to the Förster resonance. It shows three peaks because, in contrast to the data of Fig. 2 of the main text, we do not postselect the data on the destructive measurement of the data atom in $|54C\rangle$. Thus, the middle peak, which corresponds to the excitation of $|45S\rangle$ for an isolated ancilla atom, results from a failed preparation of $|54C\rangle$ for the data atom. The top and bottom peaks correspond to the excitation of the interaction eigenstates, which read

$$|\pm\rangle_{\Delta_p} = \cos(\theta)|54C, 45S\rangle \pm \sin(\theta)|53C, 45P\rangle. \quad (\text{B1})$$

In Eq. (B1), $\tan(2\theta) = \Omega/\Delta_p$, with $h\Delta_p$ the electric-field-dependent energy difference between $|53C, 45P\rangle$ and $|54C, 45S\rangle$. The frequency difference Δ_p can be expressed as a function of the electric field, through the differential Stark effect coefficients, $\alpha_{SP}^{(1)} = -3.667 \text{ MHz/(V/cm)}$, $\alpha_{SP}^{(2)} = -55.568 \text{ MHz/(V/cm)}^2$,

and $\alpha_{CC}^{(2)} = -0.339 \text{ MHz}/(\text{V}/\text{cm})^2$, between $|45S\rangle$ and $|45P\rangle$ and between $|53C\rangle$ and $|54C\rangle$, respectively:

$$\Delta_p = \alpha_{SP}^{(1)}(F - F_{\text{res}}) + (\alpha_{SP}^{(2)} - \alpha_{CC}^{(2)})(F^2 - F_{\text{res}}^2). \quad (\text{B2})$$

The energies of $|\pm\rangle_{\Delta_p}$ read

$$E_{\pm} = \frac{h\Delta_p}{2} \pm \frac{h}{2} \sqrt{\Omega^2 + \Delta_p^2}. \quad (\text{B3})$$

We analyze the spectra plotted in Fig. 2(b) of the main text and in Fig. 8 in a similar way. The central peak is well fitted by a coherent excitation profile that takes the form of

$$\mathcal{P}(\Delta) = \mathcal{P}_0 + A \frac{\Omega^2}{\Omega^2 + \Delta^2} \sin^2\left(\pi \sqrt{\Omega^2 + \Delta^2} T\right), \quad (\text{B4})$$

where Ω is the Rabi frequency of the optical excitation, $T = 0.4 \mu\text{s}$ (Fig. 2) or $0.3 \mu\text{s}$ (Fig. 8) the duration of the pulse, and the coefficients $\mathcal{P}_0$ and A account for a Δ -independent limited recapture efficiency of the ancillae and a limited efficiency of the Rydberg excitation, respectively. For the data of Fig. 2, we find $\Omega = 1.30(5) \text{ MHz}$, $\mathcal{P}_0 = 0.997(5)$, $A = 0.99(1)$. The full width at half maximum (FWHM) for this peak is $2.08(8) \text{ MHz}$.

The two outer peaks, however, are broadened by the thermal spread of the positions of the ancilla and data atoms in their traps. Thus, we rather fit Gaussian profiles to these two peaks, and find, for the data of Fig. 2, a FWHM of $4.0(2) \text{ MHz}$ for the right peak and $3.8(1) \text{ MHz}$ for the left peak. The measured broadening with respect to the central peak is compatible with a thermal spread of the atoms with a temperature of $T = 4.5 \mu\text{K}$ and a trap oscillation frequency of 20 kHz .

We plot in Fig. 8(c) the frequencies of the three observed peaks as a function of F , obtained by fitting to the data of Fig. 8(b) a sum of two Gaussian peaks and of a coherent excitation profile [Eq. (B4)]. The frequency of the central peak depends on F through the Stark effect of $|45S\rangle$ only, and can be expressed as $\nu_{\text{res}} + \Delta_0(F)$, where ν_{res} is the frequency of the $|1\rangle \rightarrow |45S\rangle$ transition when $F = F_{\text{res}}$ or, similarly, of the $|54C, 5S\rangle \rightarrow |54C, 45S\rangle$ transition without dipole-dipole interaction. The frequency difference $\Delta_0(F)$ reads

$$\Delta_0(F) = \alpha_{45S}^{(1)}(F - F_{\text{res}}) + \alpha_{45S}^{(2)}(F^2 - F_{\text{res}}^2), \quad (\text{B5})$$

where $\alpha_{45S}^{(2)} = 0.537 \text{ MHz}/(\text{V}/\text{cm})^2$ and $\alpha_{45S}^{(1)} = -12.322 \text{ MHz}/(\text{V}/\text{cm})$ are the quadratic and linear Stark coefficients of $|45S\rangle$ calculated for a field $F \approx 1.6 \text{ V cm}^{-1}$ [57]. The electric-field-dependent frequencies of the $|54C, 5S\rangle \rightarrow |\pm\rangle_{\Delta_p}$ transitions, corresponding to the two outer peaks, read $\nu_{\text{res}} + \Delta_{\pm}(F)$, with

$$\begin{aligned} \Delta_{\pm}(F) &= \Delta_0(F) + \frac{1}{2} \Delta_p(F) + \frac{E_{\pm}(F)}{2} \\ &= \Delta_0(F) + \frac{1}{2} \Delta_p(F) \pm \frac{1}{2} \sqrt{\Omega^2 + \Delta_p(F)^2}. \end{aligned} \quad (\text{B6})$$

We fit Eqs. (B5) and (B6) to the data of Fig. 8(c) with Ω , F_{res} and ν_{res} as free parameters. We find $F_{\text{res}} = 1.592(1) \text{ V}/\text{cm}$ and $\Omega = 18.2(1) \text{ MHz}$, which corresponds to an interatomic distance of $6.27(2) \mu\text{m}$. The discrepancy with the measured $6.1\text{-}\mu\text{m}$ distance with the data of Fig. 2 can be explained by the fact that the phase masks used for the two experiments, performed a couple months apart, are slightly different.

APPENDIX C: ESTIMATION OF THE ABM FIDELITIES

In Appendix C, we described the experimental imperfection model and its fits to the data.

1. Procedure

In this appendix, we describe the model used to obtain a precise estimation of the nondestructive measurement efficiencies, $\eta_{\uparrow}$ and $\eta_{\downarrow}$, defined as the probability to block the optical excitation of the ancilla to $|R_a\rangle$ when the data atom is in $|\uparrow\rangle$ and the probability to excite the ancilla to $|R_a\rangle$ when the data atom is in $|\downarrow\rangle$. The measured spectra of Fig. 2 of the main text and the values of $\mathcal{P}_A$ at resonance would be sufficient to estimate $\eta_{\uparrow}$ and $\eta_{\downarrow}$, would the DM be error-free. For a proper estimation of the ABM efficiency, we need to consider the limitations of the DM scheme. The latter are, e.g., responsible for the small peak around $\Delta = 0$, in the data of Fig. 2(b), when the data atom has been measured by the DM scheme in $|\uparrow\rangle$ (see main text).

To do so, we repeat the experiments presented in Fig. 3(a) of the main text, with and without the MW π pulse that transfers $|\uparrow\rangle$ to $|\downarrow\rangle$ after the ABM sequence; the destructive measurement effectively measures the probability to have the data atom in $|\uparrow\rangle$ and $|\downarrow\rangle$, respectively. In Figs. 9(b) and 9(c), we plot the joint probabilities to detect the data and ancilla atoms without and with the final MW π pulse. The possible outcomes are labeled with the following terminology. Filled and open symbols indicate an atom eventually detected and absent, respectively, in the final fluorescence image. The left triangles ($\blacktriangle$, $\triangleright$) correspond to the ancilla to the left of the data atom in Fig. 6, the right triangles ($\blacktriangleleft$, $\triangleleft$) to the right ancilla, and the central circles ($\bullet$, $\circ$) to the data atom. The large circle ($\bigcirc$) is specifically used to represent a data atom in $|\uparrow\rangle$. A symbol $\times$ in place of the previously defined symbols corresponds to any possible state.

Every triplet of symbols denotes the conjoined presence or absence of the three atoms. When present, an index indicates the index (1, 2, or 3) of the fluorescence image used to measure the probability of the considered event.

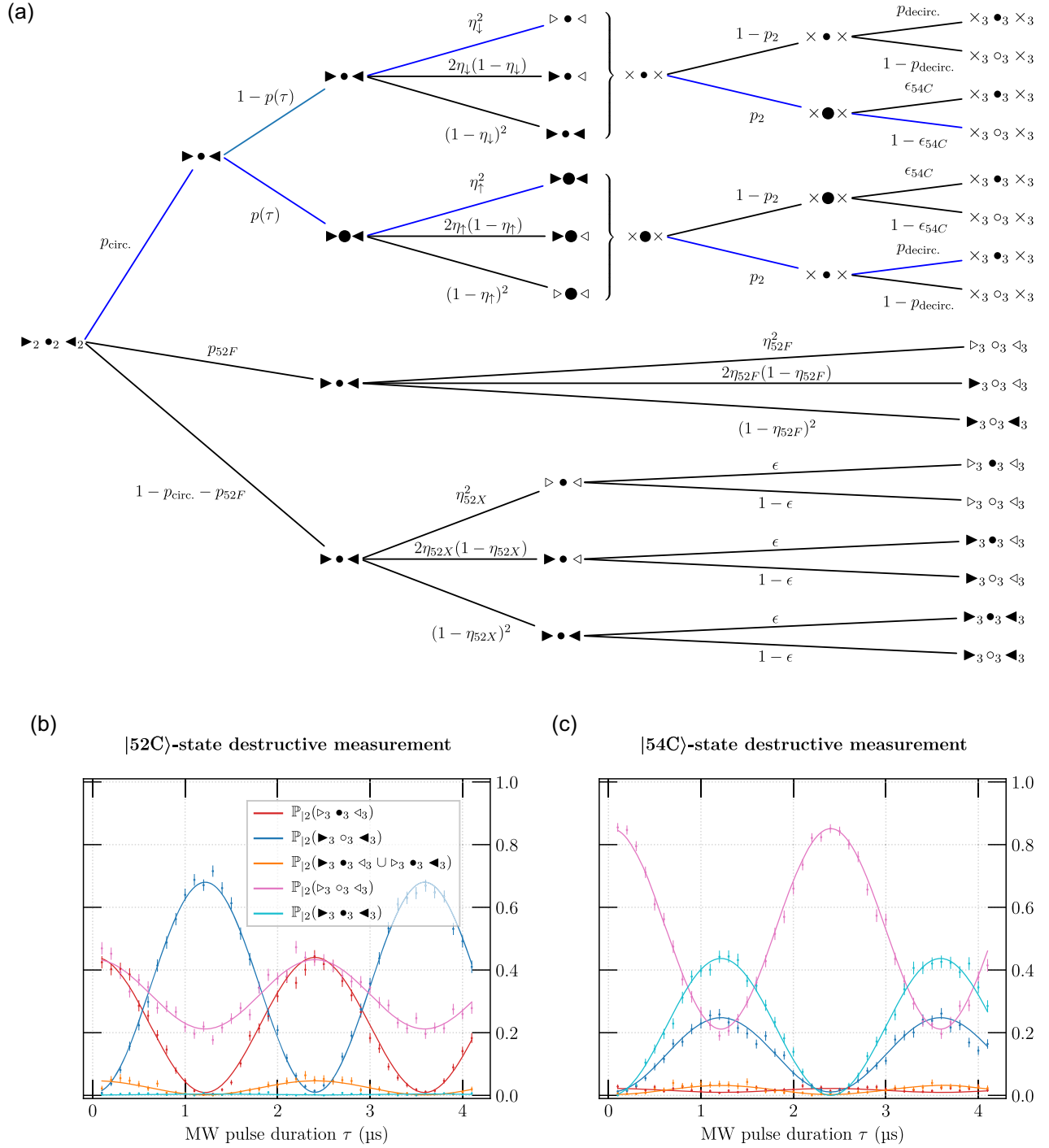


FIG. 9. Error model and fit to the data. (a) Probability tree corresponding to the measurements of Rabi oscillations with the ABM and the DM. The tree is drawn in the situation of Figs. 3(a) and 3(b) of the main text, namely, when an MW π pulse is used after the ABM sequence to measure the state $|\uparrow\rangle = |54C\rangle$ with the DM. The blue branches represent the two error-free paths. The symbol $\times$ is used for compactness: It denotes any of the ancilla symbols ($\blacktriangleright, \triangleright, \blacktriangleleft, \triangleleft$) of a previous branch. The branches drawn after the curly brackets are meant to be reproduced 3 times with the three considered configurations of ancillae. For the events with different outcomes for the two ancillae, we draw only one configuration to account for the two possible symmetric configurations. (b),(c) Measured conditional probabilities of the five relevant events without (b) and with (c) the final MW π pulse that transfers atom in $|\uparrow\rangle$ to $|\downarrow\rangle$ before the DM scheme. In the two panels, the points are measurement outcomes, with the color code given in the inset of (b). The solid lines are the model curves fitted to the data with the parameters of Table I.

The starting event $\blacktriangleright_2 \bullet_2 \blacktriangleleft_2$ corresponds to the three atoms present in the second fluorescence image, made after the rearrangement of the initially loaded array. The notation $\mathbb{P}_2(X)$ corresponds to the probability of event X conditioned on the effective realization of $\blacktriangleright_2 \bullet_2 \blacktriangleleft_2$, namely, of the preparation of the defect-free target array.

In the model we develop, the two ancillae are perfectly symmetric: The model does not distinguish the $\triangleright \times \blacktriangleleft$ and $\blacktriangleright \times \triangleleft$ events. Therefore, for these asymmetrical events, we calculate instead the sum of their probabilities, leaving only five independent triplets of events: $\blacktriangleright \bullet \blacktriangleleft$, $\blacktriangleright \bullet \triangleleft$, $\triangleright \bullet \blacktriangleleft$, $\triangleright \bullet \triangleleft$, $(\triangleright \bullet \blacktriangleleft) \cup (\blacktriangleright \bullet \triangleleft)$. We deduce the probability of the remaining event $(\triangleright \bullet \triangleleft) \cup (\blacktriangleright \bullet \blacktriangleleft)$ from the five others. We measure these five independent probabilities at the end of the experimental sequence and plot them in Figs. 9(b) and 9(c) as a function of the MW pulse duration τ .

2. Details of the model

In order to understand the contrast of the oscillations observed in Fig. 9, we develop a model summarized in the form of a probability tree sketched on Fig. 9(a). It includes 13 independent parameters, listed below.

- (i) $\eta_{\uparrow}$, measurement efficiency of $|\uparrow\rangle$.
- (ii) $\eta_{\downarrow}$, measurement efficiency of $|\downarrow\rangle$.
- (iii) $p_{\text{circ.}}$, the circularization efficiency, namely, the probability for an atom in $|52F\rangle$ to be transferred to $|52C\rangle$ with the first rf adiabatic passage.
- (iv) $p_{\text{decirc.}}$, the decircularization efficiency, namely, the probability for an atom in $|52C\rangle$ to be transferred to $|52F\rangle$ with the second rf adiabatic passage.
- (v) p_{52F} , the probability that the data atom remains in $|52F\rangle$ after the circularization process. It differs from $1 - p_{\text{circ.}}$ because the atom may be left by the circularization process in a state different from $|52F\rangle$ or $|52C\rangle$. In the following, we denote $|52X\rangle$ any state other than $|52F\rangle$, $|52C\rangle$, and $|54C\rangle$.
- (vi) p_1 , the efficiency of the first MW pulse that transfers $|52C\rangle$ to $|54C\rangle$. More precisely, the transfer probability reads, as a function of the variable duration τ of the MW pulse, $p(\tau) = p_1 \sin^2[(\pi/2)(\tau - \delta\tau)/\tau_\pi]$.
- (vii) τ_π , the duration of the aforementioned π pulse.
- (viii) $\delta\tau$, a phase parameter used to account phenomenologically for imperfect square pulses.
- (ix) p_2 , the efficiency of the second fixed-duration MW π pulse that follows the ABM sequence, used only for the data of Fig. 9(c).
- (x) ϵ_{54C} , the probability that the decircularization process transfers an atom in $|54C\rangle$ to a low- ℓ state that eventually decays to the ground state and is recaptured in the Gaussian tweezers at the end of the sequence.
- (xi) η_{52F} , the probability to excite an ancilla atom to $|R_a\rangle$ when the data atom is in $|52F\rangle$. It accounts for the fact that the static electric dipole of $|52F\rangle$ in F_{res} is strong enough to shift the excitation line of the

ancilla to $|R_a\rangle$ by 0.86(4) MHz and, therefore, to partially block the excitation.

- (xii) η_{52X} , the probability to excite an ancilla atom to $|R_a\rangle$ when the data atom is in a state $|52X\rangle$.
- (xiii) ϵ , the probability to recapture in the Gaussian tweezers at the end of the sequence the data atom in a state $|52X\rangle$.

The probabilities $p_{52F} = 3.9\%$ and $\eta_{52F} = 46\%$ have been measured independently with dedicated experiments. The 11 other parameters are free parameters of the model. Note that we single out the state $|52F\rangle$, and do not include it in the $|52X\rangle$ states, for two reasons. On the one hand, $|52F\rangle$ can partially block the excitation of the ancilla to $|R_a\rangle$, contrarily to, e.g., $|52D\rangle$, $|53C\rangle$, or other high- ℓ Rydberg levels. On the other hand, if the data atom is in $|52F\rangle$ during the ABM sequence, there is a high chance that it will be transferred to $|52C\rangle$ by the second rf adiabatic passage. For simplicity, we consider that those atoms cannot be eventually recaptured. Other Rydberg levels, included in $|52X\rangle$, may have a non-negligible probability to be recaptured, accounted for by the average coefficient ϵ .

3. Fit results

Using the probability tree, we calculate as a function of τ and of the parameters defined above the probabilities of the 5 events plotted in Fig. 9, in the two realizations of Figs. 9(b) and 9(c). The 10 probabilities are sinusoidal functions of τ with common period τ_π and phase $\delta\tau$. We express their 10 contrasts and mean values as a function of the 9 other free parameters, p_{52F} and η_{52F} . We plot the fitted curves in Figs. 9(b) and 9(c) as solid lines. We find a good agreement with the experiment for the values of the parameters given in Table I.

4. Imperfections analysis

In this section, we identify the main imperfections that limit the fitted values of $\eta_{\uparrow}$ and $\eta_{\downarrow}$. Table II sums up these different imperfections and provides estimates of their contributions to the measurement efficiencies. These contributions have been calculated or measured independently via auxiliary experiments. They have not been included in the model to avoid significant complication of the probability tree.

The measurement efficiency $\eta_{\uparrow}$ of $|54C\rangle$ is limited by losses of an ancilla atom, when it should have been shelved in $|0\rangle$ during the ABM sequence and eventually detected in the final fluorescence image. The first loss source stems from collisions with the background gas, and is estimated to 1.5% over the ~ 84 -ms delay time between the second and third fluorescence images of the experimental sequence, given the 5.8-s lifetime of atoms in the traps. The second loss source originates from photoionization of the ancillae during the excitation of the data atom to $|52D\rangle$. Photoionization is induced by the 420-nm-wavelength blue laser in the presence

TABLE I. Fitted values of the 11 free parameters of the model and of p_{52F} and η_{52F} estimated with independent measurements. The fitted probabilities are bounded to the $[0, 1]$ interval.

Parameters	Fitted value
p_{circ}	72.3(4)%
ϵ_{54C}	0.7(1)%
p_{decirc}	66.0(4)%
ϵ	4.4(4)%
$\eta_{\uparrow}$	96.6(2)%
p_1	100%
$\eta_{\downarrow}$	95.0(3)%
p_2	97.9(4)%
η_{52X}	93.7(9)%
$\delta\tau_p$	0.06(1) μs
τ_{π}	1.187(3) μs
Parameters	Measured value
p_{52F}	3.7(4)%
η_{52F}	46(2)%

TABLE II. Experiment imperfections limiting the QND measurement efficiencies.

54C>-state measurement		52C>-state measurement	
Imperfection	Estimation	Imperfection	Estimation
Collisions	1.5%	Lifetime	1.3%
Photoionization	1.4%		
Imaging and pumping	0.5%	Phase noise	0.9%
		Trap depumping	0.5%
		Optical pumping	0.4%
Total	3.4%	Total	4.1%
Measured: $1 - \eta_{54C}$	3.4(2)%	Measured: $1 - \eta_{52C}$	5.0(3)%

of the 821-nm-wavelength 1-mW-power trapping light. Additional losses are induced by the fluorescence imaging itself and the optical pumping stage (0.5%). Remarkably, these imperfections suffice to explain the value of $\eta_{\uparrow}$. In particular, we find no limitations due to an eventual imperfect dipole blockade of the excitation of the ancillae to $|R_a\rangle$.

The measurement efficiency $\eta_{\downarrow}$ of the state $|52C\rangle$ is limited by four main imperfections, that lead to a final recapture of the ancilla atom even though it should have been excited to $|R_a\rangle$ and eventually expelled from the traps. The first one is the 51- μs lifetime of $|R_a\rangle = |45S_{1/2}, m_J = +1/2\rangle$. An ancilla in $|R_a\rangle$ can decay back to $|0\rangle = |5S_{1/2}, F = 1\rangle$ prior to the second Rydberg excitation pulse that transfers $|R_a\rangle$ back to $|1\rangle$. The second limitation corresponds to an imperfect excitation of $|1\rangle$ to $|R_a\rangle$, which may stem from laser phase noise. The excitation efficiency is also limited by a limited preparation efficiency of $|1\rangle$. In particular, the 1.0-mW optical tweezers

that trap the ancillae can pump the atoms out of the $|1\rangle$ state over the 6 ms before the Rydberg excitation of the ancillae. The finite efficiency of the initial optical pumping itself to $|1\rangle$ constitutes the fourth considered imperfection.

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