Long-lived Bell states in an array of optical clock qubits

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The generation of long-lived entanglement in optical atomic clocks is one of the main goals of quantum metrology. Arrays of neutral atoms, where Rydberg-based interactions may generate entanglement between individually controlled and resolved atoms, constitute a promising quantum platform to achieve this. Here we leverage the programmable state preparation afforded by optical tweezers and the efficient strong confinement of a three-dimensional optical lattice to prepare an ensemble of strontium-atom pairs in their motional ground state. We engineer global single-qubit gates on the optical clock transition and two-qubit entangling gates via adiabatic Rydberg dressing, enabling the generation of Bell states with a state-preparation-and-measurement-corrected fidelity of 92.8(2.0)% (87.1(1.6)% without state-preparation-and-measurement correction). For use in quantum metrology, it is furthermore critical that the resulting entanglement be long lived; we find that the coherence of the Bell state has a lifetime of 4.2(6) s via parity correlations and simultaneous comparisons between entangled and unentangled ensembles. Such long-lived Bell states can be useful for enhancing metrological stability and bandwidth. In the future, atomic rearrangement will enable the implementation of many-qubit gates and cluster state generation, as well as explorations of the transverse field Ising model.

A s the essential resource in quantum science, quantum entanglement enables a broad set of applications in computing, cryptography, and materials science, to name a few. One powerful application arises in metrology, where greater sensitivity and higher-bandwidth sensors are afforded by the properties of entangled multiparticle quantum states1–8. Combining such enhancements with state-of-the-art time and frequency metrology9–14—namely, optical atomic clocks—has been a defining goal in this field of quantum metrology; the construction of a quantum-enhanced optical clock has broad implications for geodesy15,16, gravitational-wave detection17–19 and the search for physics beyond the standard model20.

A variety of approaches exist for creating metrologically useful entanglement. In neutral-atom optical lattice clocks, a number of methods have been proposed using cavity quantum electrodynamics, Rydberg interactions, or collisional interactions21–26—indeed, recently, spin-squeezed states on an optical clock transition have been generated using collective cavity quantum electrodynamics interactions27. In trapped ions, proposals and implementations of entanglement on optically separated qubits rely on spin–spin interactions mediated by Coulombic crystal modes, allowing efficient entanglement generation and Greenberger–Horne–Zeilinger states with as many as 24-ion optical qubits28 or photonic quantum networking between spatially distributed single-ion clocks29. Given the control and measurements possible in modern atomic clock architectures, these systems offer the possibility of merging quantum information concepts with precision measurement. Utilizing entangling gates and protocols from quantum information science broadens the set of realizable quantum states, enabling new protocols for quantum metrology, while quantum error-correcting codes have opened routes for enhanced quantum sensing in the presence of noise and imperfect quantum resources30–34. In the context of optical atomic clocks, a number of recent protocols have advanced this union as a route to quantum-enhanced measurements30,31,33, and in a very recent demonstration, the variational optimization of phase sensitivity at optical frequencies was implemented on a 26-ion quantum processor32.

This theoretical and experimental momentum emphasizes the promise of combining control with scalability, to produce large-scale, entangled clocks amenable to versatile control and measurement schemes31,36,37. An outstanding challenge in this regard is to pair these capabilities with long-lived quantum coherence so that many-particle states can be fruitfully leveraged for quantum-enhanced measurements at long interrogation times. A promising candidate for realizing this synergy is the recently demonstrated tweezer clock, which not only supplies long-lived atomic coherence and high stability34–40 but also large atom number39–42, high-fidelity Rydberg interactions43–46 and microscopic control to engineer precisely tailored quantum systems.

In our previous work, we demonstrated state-of-the-art relative stability and atomic coherence using a tweezer-array clock38,40. Here we advance our atomic control of the clock transition to elevate it into a high-fidelity qubit. We then implement entangling gates between these qubits through Rydberg excitations out of the excited qubit state, which realizes controllable Ising-type interactions41–44.

With these capabilities, we achieve high-fidelity long-lived entanglement in an ensemble of Bell state pairs (Fig. 1a)45,46. Here we employ an adiabatically resonant coupling of the qubit excited state to a Rydberg state to implement a controlled phase gate between two atomic qubits, up to single-atom lightshifts45. Setting the controlled phase of our gate to be $\pi/2$ yields a controlled-$\sqrt{Z}$ gate and enables the production of the Bell state $|\psi\rangle = \frac{1}{\sqrt{2}}(|gg\rangle + i|ee\rangle)$, which has metrological relevance because the superposition components have a large energy separation. As such, clocks based on

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such states accrue phase more quickly than clocks based on unentangled atoms (Fig. 1b,c,d). We determine the fidelity with which we produce Bell states to be $F = 92.8(2.0)\%$ over a 6 x 8 array of atom pairs. To better understand near-term prospects for quantum metrological enhancement, we study the lifetime of these states in several ways. We observe the parity oscillation contrast and resulting fidelity to decay with a Gaussian $1/e$ time constant of $\sigma_L = 407(13) \text{ ms}$. However, the loss of atom–laser coherence can occur without the loss of atom–atom coherence, or in this case, relative coherence between distinct Bell states. Taking inspiration from the measurements of atomic coherence that reject the relative phase noise of the interrogation laser$^{14,40,56}$, we measure Bell state parity and single-atom $S_e$ correlations to infer a $1/e$ exponential Bell state coherence time of $\tau_c = 4.2(6) \text{ s}$.

Our experiments begin with a tweezer-defined doublet array stochastically filled with $^{88}\text{Sr}$ atoms (Fig. 1a). The atoms are implanted into a single plane of a co-located three-dimensional (3D) optical lattice, imaged, and subsequently cooled into their 3D motional ground state to achieve high-fidelity atomic control (Methods).

We then apply a sequence of global laser pulses that drive the $^{1}S_0$ (|g⟩) $\leftrightarrow$ $^{1}P_1$ (|e⟩) clock transition resonantly with Rabi frequency $\Omega$, as well as the |e⟩ $\leftrightarrow$ |r⟩ Rydberg transition (to $5s40d^1D_2$; $m=0$) with Rabi frequency $\Omega_r$ and positive detuning $\Delta$ (Fig. 2a). Finally, we image the excited clock state atoms.

We implement a controlled phase gate for two clock-transition qubits using adiabatic Rydberg pulses, following the proposal of Mitra, et al.$^{54}$. Each Rydberg pulse ramps $\Omega$ and $\Delta$ so that the clock state atoms adiabatically follow the instantaneous dressed clock-like eigenstate in the presence of the Rydberg coupling laser. This lightshifts the energy of a single clock state atom by $E_{\text{LS}}^{(1)} = \frac{\hbar}{2} \left( -\Delta + \sqrt{\Delta^2 + \Omega^2} \right)$, where $\hbar$ is the reduced Planck's constant. The energy shift of two non-interacting clock state atoms is $2E_{\text{LS}}^{(1)}$. As shown in Fig. 2b, at close distances, such as two atoms within a single doublet, Rydberg blockade prevents the excitation of both atoms to the Rydberg state. The Hilbert space of two clock state atoms then reduces to an effective two-level system with states $|ee\rangle$ and $|b\rangle = \frac{1}{\sqrt{2}} (|ee\rangle + |re\rangle)$ driven by a collectively enhanced Rabi frequency $\sqrt{2}\Omega$, and detuning $\Delta$. This lightshifts the energy of the $|ee\rangle$ state by $E_{\text{LS}}^{(2)} = \frac{\hbar}{4} \left( -\Delta + \sqrt{2\Delta^2 + \Omega^2} \right)$. The difference $\kappa = E_{\text{LS}}^{(2)} - 2E_{\text{LS}}^{(1)}$ is called the entangling energy, as it sets the energy scale of entangling operations$^{54}$.

This adiabatic Rydberg pulse implements a controlled phase gate on the two-qubit states of a filled doublet, once the single-atom phase is removed. That is, the Rydberg pulse realizes the Hamiltonian

$$H_{\text{r}} = E_{\text{LS}}^{(1)} S_z + \kappa |ee\rangle \langle ee|,$$  

where $S_z = \mathbb{1} \otimes \sigma_z/2 + \sigma_z/2 \otimes \mathbb{1}$, $\sigma_z$ is the Pauli $z$ matrix and $\mathbb{1}$ is the identity matrix acting on the $(|e\rangle, |g\rangle)$ basis. We only consider single- and two-particle effects because the spacing between doublets is chosen to be large compared with the Rydberg interaction range (Fig. 1a). By implementing two of these Rydberg pulses separated by a $\pi$ pulse, we remove the single-atom phase from $S_z$; therefore, the full sequence can be interpreted as containing two controlled phase gates, $U_{\phi} = |gg\rangle \langle gg| + |ee\rangle \langle ee| + |ge\rangle \langle eg| + e^{i\phi} |eg\rangle \langle ce| + |ce\rangle \langle ce|$, with phase $\phi = -\frac{\sqrt{2}\Delta}{\Omega}$ (Fig. 2c).

Together, these three pulses implement Ising-type interactions $H_{\text{eff}} = \kappa S_z^2$. We note that although fixed, large-detuning Rydberg dressing would produce equivalent coherent
dynamics, adiabatically following the instantaneous eigenstates to resonance increases the ratio of interactions to dissipation, improving the gate fidelity (Supplementary Information).

We reveal the dynamics of the controlled phase gates by including them inside a Ramsey interferometry sequence (Fig. 2d, top), where the gate appears between two \(\pi/2\) pulses. The first pulse establishes a superposition with a well-defined phase between \(|g\rangle\) and \(|e\rangle\) for each atom. The spin-echoed Rydberg pulses introduce an additional phase \(\phi\) for the \(|ee\rangle\) and \(|gg\rangle\) components, and the final \(\pi/2\) pulse closes the interferometer, converting those phases into an oscillation between \(|gg\rangle\) and \(|ee\rangle\) removing the population from \(|gg\rangle\) and \(|ee\rangle\). The two \(\pi/2\) pulses convert the effective Ising interactions of the middle three pulses into an \(S_z^r\) interaction, which rotates \(|gg\rangle\) into \(|ee\rangle\). As shown in Fig. 2d, for short gate times, we observe these dynamics via \(S_z^r \equiv \langle S_z^r \rangle\) and \((P_{gg} + P_{ee})\) population measurements in filled doublets. For longer gate times, we observe that the failure of the spin echo causes a decay of \(S_z^r\) oscillations, as revealed in the dynamics of the single-atom \(S_z\) relaxing to zero. Furthermore, this causes the population to leave the \(|gg\rangle\)–\(|ee\rangle\) subspace.

A zero crossing of \(S_z\) corresponds to the application of controlled phase gates of \(\phi = \pi/2\) in under a microsecond, and after the full pulse sequence (Fig. 2d), the probability of the Bell state \(|\psi\rangle = \frac{1}{\sqrt{2}}(|gg\rangle + i|ee\rangle)\). We benchmark our gates through fidelity \(F\), which we use to create these Bell states, that is, \(F = \frac{1}{2} (P_{gg} + P_{ee} + C)\) (ref. 1). This depends on two factors: the sum of populations in the relevant states \((P_{gg} + P_{ee})\) as well as the coherence between them as accessed by contrast \(C\) in a parity oscillation measurement1. At the first \(S_z\) zero crossing, we find \((P_{gg} + P_{ee}) = 0.955(28)\) or \(0.922(27)\) without state-preparation-and-measurement (SPAM) correction (Fig. 3a and Supplementary Information). To extract the Bell state coherence, we add a final ‘analysis’ \(\pi/2\) pulse about a variable axis after completing the spin-echo sequence. In Fig. 3b, we plot the measured parity \(P \equiv \langle P \rangle = \langle \delta_1^r \delta_2^r \rangle\), where the superscript ‘l’ and ‘r’ denotes the left and right sites, respectively, within a doublet and the brackets imply averaging over the array and repetitions of the experiment. The amplitude of a sinusoidal fit versus the analysis phase yields the parity oscillation contrast \(C = 0.919(28)\) and a resulting Bell-state fidelity of \(F = 92.8(2.0)\%\) (without SPAM correction, \(C = 0.838(19)\) and \(F = 87.1(1.6)\%\). The reported fidelity is reduced by 0.94(4)% to account for a systematic overestimation of \(P_{gg}\) due to Rydberg state loss (Supplementary Information), and the uncertainty includes population and contrast uncertainty as well as uncertainty from SPAM correction, primarily from fluctuations of the imaging loss at the percent level.

The adiabatic Rydberg ramp speed and detuning endpoints are optimized to satisfy several competing constraints. Operating faster reduces the loss from the finite Rydberg-state lifetime (10.9(4) \(\mu\)s (Supplementary Information)), whereas operating slower results in greater adiabaticity so that atoms more completely return to the qubit states (Supplementary Information). We balance these competing effects with ramp durations of 250–350 ns (Methods). Master equation modelling (Fig. 2c and Supplementary Information) guided the experimental optimization of these parameters (Methods).

For metrological use, a long coherence time for the entangled state is desired. In Fig. 4, we investigate this timescale in several ways. First, we can simply wait before measuring the populations...
and parity oscillation contrast. Although the populations decay slowly (12.3(1.0) s exponential time constant for $P_{ab}$ (Supplementary Information)), we find that the parity oscillation contrast decays rapidly, with a Gaussian 1/e constant of $\sigma_z = 407(13)$ ms (Fig. 4a). This loss of Bell-state-laser coherence could arise due to laser noise (in the rotating frame of the atoms), rather than atomic dephasing or decoherence effects. To investigate this possibility, we prepare the ensemble of single atoms in the array into a superposition state $\hat{\sigma}_z = (\hat{\sigma}_z - 1|\hat{\sigma}_z|)$ while preserving the Bell states’ parity via an additional $\pi/2$ pulse about a variable axis is implemented by the composition of a single variable gate $\hat{g}_g$ and a static $\pi/2$ pulse and produces two-particle states that oscillate in parity $\hat{P}_g = \hat{g}_g \hat{\sigma}_z \hat{g}_g^\dagger$. The resulting data (blue points, SPAM corrected; error bars are s.e.m. from averaging over the array and 15 repetitions of the experiment) is fit by a sinusoid (blue curve; fit error band represents 1σ uncertainty). The amplitude of the sinusoidal fit yields $C = 0.919(28)$. The uncertainty arises in approximately equal parts from the fit-parameter estimation uncertainty and systematic imaging loss fluctuations.

The parametric plots of the entangled ensemble parity against the entangled ensemble parity (Supplementary Information). Importantly, these correlations do not, by themselves, certify the presence of a Bell state, since there are non-entangled states that could result in similar signals. However, Fig. 3 provides tomographic evidence for the initial generation of a Bell state, and these correlations then provide the Bell state coherence time with assumptions that certain exotic decay processes do not occur (Supplementary Information). The comparison of the ensemble of Bell states to the ensemble of single atoms (Fig. 4c) is a simultaneous comparison between an entangled and unentangled pair of optical atomic clocks. The parametric plots of the entangled ensemble parity against the unentangled ensemble parity and $S_z$ at hold times of less than a second form 2:1 and 1:1 Lissajous figures with zero relative phase, that is, a parabola and a line, respectively (Fig. 4d,e). The operation of a quantum-enhanced optical atomic clock requires the estimation
of the relative phase difference between the reference and target clocks. To examine this, we inject a relative phase difference of \(\pi/2\) with a third controlled-\(\sqrt{Z}\) pulse. This opens the 2:1 and 1:1 Lissajous figures into a figure of eight and ellipse, respectively, with a relative phase of 0.94(7) radians measured via ellipse fitting (Fig. 4f,g). Future operation of this system in a quantum-enhanced simultaneous clock comparison may require establishing an optimal initial-phase difference and subsequently estimating the opening angle of the figure of eight or ellipse at an optimal interrogation time, whereas in situ comparison of an entangled and unentangled ensemble opens new routes for characterizing clock systematics and retaining enhanced metrological sensitivity in the presence of laser noise\(^{31,36,37}\).

The two measurements of the Bell state coherence time from correlations are mutually consistent and may be combined to yield \(\tau_0 = 4.2(6)\) s. This, combined with our understanding of single-atom and laser decoherence mechanisms, provides an expected exponential decay time for the parity oscillation contrast of 0.69(10) s, longer than that shown in Fig. 4a. Nevertheless, the Bell state coherence time from correlations may be compared with an experimentally determined expectation given by half the single-atom atomic coherence time, that is, the decay rate of the \(S_x - S_y\) correlation of single atoms, namely, \(\tau_{0,\text{SA}} = 12.2(5)\) s. This value is understood in terms of single-atom population loss, decay and decoherence processes, up to an unexplained single-atom lossless decoherence process with a four-minute timescale. This value is also consistent with the observed Bell state population dynamics, indicating the absence of correlated loss or decay processes (Supplementary Information). The nearly factor-of-three reduction in the Bell state coherence time from in situ correlation measurements could indicate the existence of an unidentified additional Bell state decoherence process with an approximately six-second timescale. However, the late-time behaviour of the correlations is consistent with single-atom measurements, and we hypothesize that the observed discrepancy is the result of technical noise that results in an additional, faster decoherence timescale at early times (Supplementary Information).

In this work, we have demonstrated the parallel generation of entangled Bell states on an optical clock transition using a novel spin-echoed adiabatically resonant Rydberg pulse sequence. We characterize the SPAM-corrected Bell-state fidelity to be \(F = 92.8(2.0)\%\) and measure a 4.2(6)s array-averaged coherence time. These results establish a firm foundation for future explorations of quantum-enhanced optical frequency metrology at the stability frontier. A crucial near-term goal in this direction is to extend these results from pairs to clusters of several-to-tens of atoms, as well as to identify and generate the particular optimal quantum states of these clusters for a clock stability measurement\(^{36,37}\). Entangling operations both before and after the interrogation time may further enhance the clock stability\(^{31,36-40}\)
More broadly, we have demonstrated capabilities useful throughout quantum science. Although previous work has shown the control and detection of highly cohered qubits and scalability to large system sizes, here we have introduced maximally entangling gates, which enables a host of quantum applications. Particularly exciting goals include the generation of cluster states and with the addition of local measurements, subsequent demonstration of one-way quantum computing. Explorations of the transverse-field Ising model with finite-range interactions are directly within reach. The combination of tweezer-based atom rearrangement and a 3D optical lattice further allows access to two-dimensional (2D) Hubbard physics with finite-range interactions.

Note added in proof: during preparation of this manuscript, we became aware of related work implementing a Mølmer–Sørensen gate with caesium atoms.

Online content
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Methods

Apparatus. Loading, cooling and imaging of $^{88}\text{Sr}$ atoms is described elsewhere. The radio-frequency source for the generation of the 515 nm tweezer array is described elsewhere. We use a 3D clock-magic optical lattice at 813 nm as the primary reference potential for imaging, clock rotations and Rydberg spectroscopy (Extended Data Fig. 1 and Supplementary Information). This potential enables efficient cooling and high-contrast clock rotations, as discussed below.

We open the 2$s-1$s transition of $^{88}\text{Sr}$, optical clock transition by applying a vertically oriented magnetic field with a magnitude of up to 550 G. This is accomplished using our main (magneto-optical trap) coils driven by a remotely controlled Delta Elektronika (SM30-200) supply. We drive the clock transition using approximately 13 mW of 698 nm laser light focused by a 150 mm cylindrical lens. The clock laser source is a series of three diodes injection locked to the light stabilized to a cryogenic silicon cavity. The path from the reference laser to our experiment includes 60 m of fibre and approximately 2 m of free space, which are noise cancelled, with the reference mirror attached to the main objective mount. In addition, there is approximately 1 m of propagation that includes the second injection-locked diode which is not noise-cancelled. Given the stability shown in our previous work with similar lengths of uncancelled free-space propagation, we do not expect the added phase noise from the uncancelled path length to limit the clock-pulse fidelity. As shown in Extended Data Fig. 2a, we can drive the clock transition with a Rabi frequency in excess of 1 kHz and with a $\pi$-pulse fidelity of 99.41(57)%. This represents a substantial increase over approximately 80% fidelity shown in our previous work.

Our high-fidelity clock rotations are enabled by the improved spatial confinement afforded by a deep, power-efficient 3D lattice as opposed to shallow tweezers. Tight confinement enables resolved sideband cooling after transfer into the lattice, yielding a reduced average phonon number $n = 0.04^{+0.15}_{-0.08}$ in the direction of clock laser propagation. Tight confinement also manifests as a high 99 kHz trap frequency and correspondingly small Lamb–Dicke parameter $\eta = 0.22$. Both lower temperature and smaller Lamb–Dicke parameter result in the reduced motional dephasing of Rabi oscillations, yielding an expected (ideal) $\pi$-pulse fidelity of 99.96$^{+0.02}_{-0.09}$. In typical experiments, however, we observe some heating that limits the clock $\pi$-pulse fidelity to a maximum of 99.79% (Supplementary Information), which may be compared with our observed clock $\pi$-pulse fidelity of 99.41(57)%.

We connect the clock state to a Rydberg state using an ultraviolet (UV) laser near 317 nm with approximately 1 W power. The UV laser system first produces 634 nm light via single-pass sum-frequency generation of 1,066 and 1,569 nm. A final stage of second-harmonic generation inside a resonant cavity converts 634 nm light into 317 nm light (Supplementary Information). We drive the transition from the clock state to the 5s40d1D (m = 0) Rydberg state with Rabi frequencies up to $\Omega = 2.8 \times 10^{18}$ Hz. We choose a D state rather than an S state due to its greater dipole matrix element, which directly increases the entangling energy in our adiabatically resonant pulse scheme, whereas the reduced van der Waals $C_6$ interaction coefficient is still sufficient to ensure blockade at the 1.1 pm interdoublet separation and enabling denser doublet spacings. We observe Rydberg spectra and Rabi oscillations through the loss of clock-state atoms (Extended Data Fig. 2b). Perhaps surprisingly, the observed contrast is consistent with the 1/9 branching ratio of intermediate triplet S states back into the clock state versus all $5s^5p^5$ states.

Experiments. Our experiments begin by preparing a magneto-optical trap, loading and cooling a single atom in a 515 nm tweezer array, transferring those atoms into a single plane of a 3D optical lattice at 813 nm, imaging the location of those atoms and cooling the atoms, resulting in >90% of the atoms in their 3D motional ground state (Extended Data Fig. 3, state prep). This process has been largely described in our previous work, except for the transfer into and cooling of the 3D optical lattice. The loss due to transfer from tweezers to well-aligned lattice sites is negligible; however, the global alignment of the tweezer array to the underlying lattice is not perfectly commensurate, resulting in ‘slips’ in the spacing of atoms in the lattice, which causes variable imaging performance. Furthermore, we find that the relative spatial phase between the tweezer array and lattice drifts by a lattice site over the course of several hours. Improving the alignment and stability of this system will be the subject of future work.

Ground-state cooling proceeds similarly to cooling in 813 nm tweezers described in previous work. Achieving a magic angle between the applied magnetic field and optical polarization enables efficient cooling to the 3D motional ground state. This does require, however, that the axial lattice be vertically polarized, which makes it less power efficient and leads to the formation of an out-of-phase horizontally polarized lattice, too. Nevertheless, we typically achieve 3D ground-state fractions above 90%, although the exact performance can vary on a day-to-day basis. Having prepared an ensemble of ground state atoms in known positions in the lattice, we are ready to perform our experiments. To open the clock transition, we apply a large magnetic field. Although we can apply a field of up to 550 G to yield a clock Rabi frequency above 1 kHz, we typically apply a 55G field, resulting in a 100 Hz clock Rabi frequency. The use of lower fields slows the gate and does not substantially affect the gate fidelity, but it does reduce the magnetic-field noise, which otherwise limits parity oscillation decay times to several tens of milliseconds due to the frequency-doubling quadrupole Zeeman shift of the clock state (Supplementary Information).

Throughout the experimental sequence, we hold the axial lattice low at approximately 8e, whereas we ramp the 2D lattice high to 420e, for all the clock pulses, low to 12e, for all the Rydberg pulses and 37e, for all the hold times. Note that e refers to the recoil energy of a single 813 nm photon. Ramping the 2D lattice high provides strong confinement for a low Lamb–Dicke parameter ($\eta = 0.22$), which enables high-fidelity clock pulses, whereas ramping the 2D lattice low reduces the total lattice-induced lightshift on the Rydberg line to ~200 kHz. This is sufficiently low such that lattice inhomogeneity does not affect gate fidelity, and leaving the lattice on provides atomic confinement. Experiments with the lattice switched off during the UV pulses can result in atomic heating, increased loss and lowered gate fidelity.

We implement the adiabatic Rydberg-dressing gate scheme of Mitra, et al. due to its reduced sensitivity to several sources of gate errors relative to the resonant gate clock of Levine, et al. In particular, we benefit from the robustness to blockade errors or Rydberg laser intensity (Rabi frequency) errors that populate the doubly excited Rydberg state. Moreover, the timing resolution of our experimental system would necessitate a reduction in the Rydberg Rabi frequency to implement the resonant gate scheme.

Finally, in the relevant experiments with a variable hold time, we perform a spin-echo clock $\pi$ pulse to remove dephasing due to spatial inhomogeneity of the clock frequency. This inhomogeneity arises from the lattices that confine the atoms—since their optical frequencies are separated by 160 MHz, both lattices cannot simultaneously be at the clock-magic wavelength. Combined with the ~50% peak-to-peak inhomogeneity of the combined lattice depth across the atom array, this yields a lightshift inhomogeneity of ~0.1 Hz. This causes a significant reduction in the various correlation decay times measured in Fig. 4. By introducing a $\pi$ pulse in the middle of the hold time, we greatly reduce the effects of inhomogeneity across the array, as evident in the spatially resolved parity–parity correlation measurements (Extended Data Fig. 4).

Data availability

The experimental data presented in this manuscript are available upon request and through Zenodo (https://doi.org/10.5281/zenodo.6626065).

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Author contributions

The experiment was designed and built by N.S., A.WY, W.E. and A.M.K. N.S., A.WY. and WJE. operated the experiment and analysed the data. All the authors contributed to the manuscript.

Competing interests

The authors declare no competing interests.

Additional information

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Extended Data Fig. 1 | FIG. M1. Trapping potentials of our apparatus. Our experiments make use of three separate optical trapping potentials. After producing a cold atomic sample via standard MOT techniques, we load atoms into an array of 515 nm optical tweezers projected through a high-numerical-aperture objective. This forms a 2d array of atoms which are then transferred into a single plane of a 3d optical lattice which is formed by two separate optical systems. First, a 2d bowtie lattice is formed using a single beam redirected by a series of mirrors and lenses. A second, axial lattice is projected from the side by directing two path-length-matched beams into a single aspheric lens.
Extended Data Fig. 2 | FIG. M2. Single atom control. a) Tight atomic confinement and a 550 G magnetic field enable a clock Rabi frequency above 1 kHz with π-pulse fidelity well above 99%; however, fidelity is limited to 99.4(57)% in part by atomic heating (which also causes dephasing at later times, see SI). Inset shows the corresponding clock transition π-pulse spectrum. b) A UV laser near 317 nm drives a clock-Rydberg transition with Rabi frequency $\Omega = 2\pi \times 13$ MHz. We detect the Rydberg state via loss of the clock state atoms. The horizontal dashed line represents an 89% Rydberg detection fidelity estimated from branching ratios of intermediate triplet $S$ states (see text).
Extended Data Fig. 3 | FIG. M3. Experimental Sequence. We present a more detailed view of the experimental sequence that includes the clock laser pulses (red), the Rydberg laser pulses (purple), the 2d lattice (orange), the axial lattice (pink), and the z-component of the magnetic field (blue). Typical lattice ramping times are 3-10 ms, and we allow the magnetic field to ramp and settle for 100 ms. Although we can achieve clock Rabi frequencies as high as 1 kHz, we operate a factor of 10 slower to reduce the effect of magnetic field noise. As such, the total gate time from the first \( \pi/2 \)-pulse through the second \( \pi/2 \)-pulse is 22.2 ms, with 12 ms arising from the four 2d lattice ramps within the gate sequence.
Extended Data Fig. 4 | FIG. M4. Removing inhomogeneous lightshifts across the atom array. The lattice potential that confines the atoms during hold times is not perfectly magic or homogeneous, which can limit array-averaged correlation signals. By increasing lattice laser detuning from the magic condition, we more clearly reveal the effects of a non-magic-lattice inhomogeneity in spatially resolved parity-parity correlation measurements at short times (**a,d**), after 1 second (**b,e**), and after 2 seconds (**c,f**). The location of each pixel corresponds to a displacement vector between fully filled doublets, while the color corresponds to the product of the parities of each doublet with that displacement vector. **a-c** Even in a shallow lattice condition, significant residual inhomogeneous lightshifts result in negative parity-parity correlations appearing along the main diagonal, reducing the array-averaged parity-parity correlator. **d-f** Introducing a π-pulse half way through the hold time spin-echoes away the inhomogeneity, increasing the array-averaged parity-parity correlator.
Extended Data Table 1 | TABLE M1. Experimental Parameters. Select parameters are tabulated for each dataset presented in the main text.

| Parameter                                                                 | Figure 2d | Figure 3a,b | Figure 4a-e | Figure 4a | 4f.g |
|---------------------------------------------------------------------------|-----------|-------------|-------------|-----------|------|
| Rydberg detuning (initial, final) (2π × MHz)                              | (40,2)    | (22,2)      | (22,5.8)    | (40,4)    | (40,8) |
| UV ramp time (ns)                                                         | 350       | 250         | 250         | 350       | 350   |
| Rydberg Rabi frequency (2π × MHz)                                        | 14        | 16          | 16          | 10        | 16    |
| Resonant gate time (ns)                                                   | varied    | 21          | 20          | 250       | 115   |
| Lattice states during UV pulses (2d, axial) (E_r)                         | (0,0)     | (12,8)      | (12,8)      | (0,0)     | (0,0) |
| Magnetic field (during hold time if applicable) (G)                       | 550       | 55          | 55 (2.8)    | 55 (55)   | 55    |