

Supplemental Material for: “Fast Nondestructive Readout for High-Clock-Rate Atom Array Quantum Processor”

I. EXPERIMENTAL SYSTEM

After loading millions of atoms into a magneto-optical trap (MOT) over 100 ms, the atoms are stochastically loaded with 3D polarization-gradient cooling into an 852 nm optical tweezer array generated by a spatial light modulator, with a trap depth of 1 mK and a spacing of 5 μm . The average loading efficiency of 52% is determined using fluorescence imaging with two global counter-propagating, π -polarized beams red-detuned from the $5S_{1/2} |F=2\rangle \leftrightarrow 5P_{3/2} |F'=3\rangle$ transition. The fluorescence is collected with a 0.55 numerical-aperture (NA) objective and imaged onto a qCMOS camera, yielding a detection fidelity of 99.97%. The trapped atoms are subsequently cooled by 1D grey molasses in the same beam geometry, reducing the average atom temperature to 14 μK . This cooling mechanism is simultaneously amenable to recooling in a finite magnetic field with a polarization-gradient configuration (detailed in Section V).

A 10 G bias magnetic field defines the quantization axis. We initialize the atoms in the stretched ground state $5S_{1/2} |F=2, m_F=2\rangle$ by optical pumping, achieving a pumping-to-depumping ratio exceeding 1000 after optimizing the optical polarization. Single-qubit control is then implemented by a global Raman drive between the two stretched hyperfine ground states $5S_{1/2} |F=1, m_F=1\rangle$ and $5S_{1/2} |F=2, m_F=2\rangle$, with a two-photon Rabi frequency of $\Omega = 2\pi \times 3.86$ MHz. The Raman beam is generated by phase modulating a laser using an electro-optic modulator and converting this phase modulation into amplitude modulation via reflecting from a chirped Bragg grating[1].

While the configuration in this work performs state readout on a 10×10 atom array, the architecture can be extended to larger arrays. The integrated 852 nm acousto-optic deflectors (AOD) equip the platform with the capability to dynamically steer atoms to positions optically conjugate to the fibre facets, enabling high-fidelity FNDR measurements within a zone-based architecture. To manage these complex experimental sequences, a custom Real-Time Messaging and Queuing (RTMQ) control system is utilized to provide low-latency execution.

II. NON-DESTRUCTIVE READOUT

Readout parameters and sequence. We demonstrate nondestructive readout by delivering two counter-propagating, σ^+ -polarized probe beams along the quantization axis, denoted P_1 and P_2 . These beams are red-detuned by $\Gamma/2$ from the free-space D2 cycling transition,

with a saturation parameter $s_0 \approx 5$. To mitigate the excited-state anti-trapping effect and tweezer-induced inhomogeneous light shifts, the tweezer and probe beams are strobed out of phase as illustrated in Fig. S1(b). Specifically, the tweezers operate at a 1MHz (0.33 μs OFF, 0.67 μs ON) with a rise/fall time of 63 ns, alongside similarly fast probe pulses. To prevent temporal overlap between the probe and tweezer light, the probe beam is pulsed for only 0.2 μs within the tweezer-OFF window. Otherwise, such overlap would shift the $F=2 \rightarrow F'=2$ transition closer to the probe frequency, exacerbating bright state leakage to dark state. During the FNDR stage, P_1 and P_2 are pulsed alternately in order to suppress the momentum-space quantum walk described by the Su-Schrieffer-Heeger (SSH) model [2]. Background noise and stray light scattering are effectively suppressed using an FPGA-controlled temporal gate, which restricts photon registration to a 0.3 μs window aligned with the active probe phase. Benefiting from this temporal filtering, alongside the high collection efficiency of the fibre array and the ultra-low dark counts of APD array, the readout achieves a signal-to-noise ratio of 144 and 80 for each side of collection, respectively.

Joint state identification. To improve the readout performance, it is essential to achieve a high photon collection efficiency. As illustrated in Fig. S1(a), we employ two confocal high-numerical-aperture objectives (NA 0.55 and 0.5, respectively)[3]. The collected photons are subsequently routed into two independent fiber-arrays [4], yielding an overall collection efficiency of 2.18%. However, the two optical paths exhibit distinct overall collection efficiencies. This asymmetry arises from variations in transmission efficiency, numerical aperture, and optical alignment between the two opposing objectives, thus rendering a direct summation of the raw counts suboptimal. To optimally combine the two detection signals, we define a weighted joint photon number, $w_1 n_1 + w_2 n_2$, where n_1 and n_2 denote the photon counts registered by fibre arrays 1 and 2, respectively. The atom is assigned to the bright state when this joint photon number exceeds a specific threshold C , establishing a linear decision boundary:

$$w_1 n_1 + w_2 n_2 > C,$$

The optimal linear discriminant can be derived from maximum-likelihood estimation (MLE). Assuming Poisson distribution for the bright (B) and dark (D) states with mean photon numbers $\mu_{S,1}$ and $\mu_{S,2}$, the joint likelihood is

$$P(n_1, n_2 | S) = \prod_{i=1}^2 \frac{\mu_{S,i}^{n_i} e^{-\mu_{S,i}}}{n_i!},$$

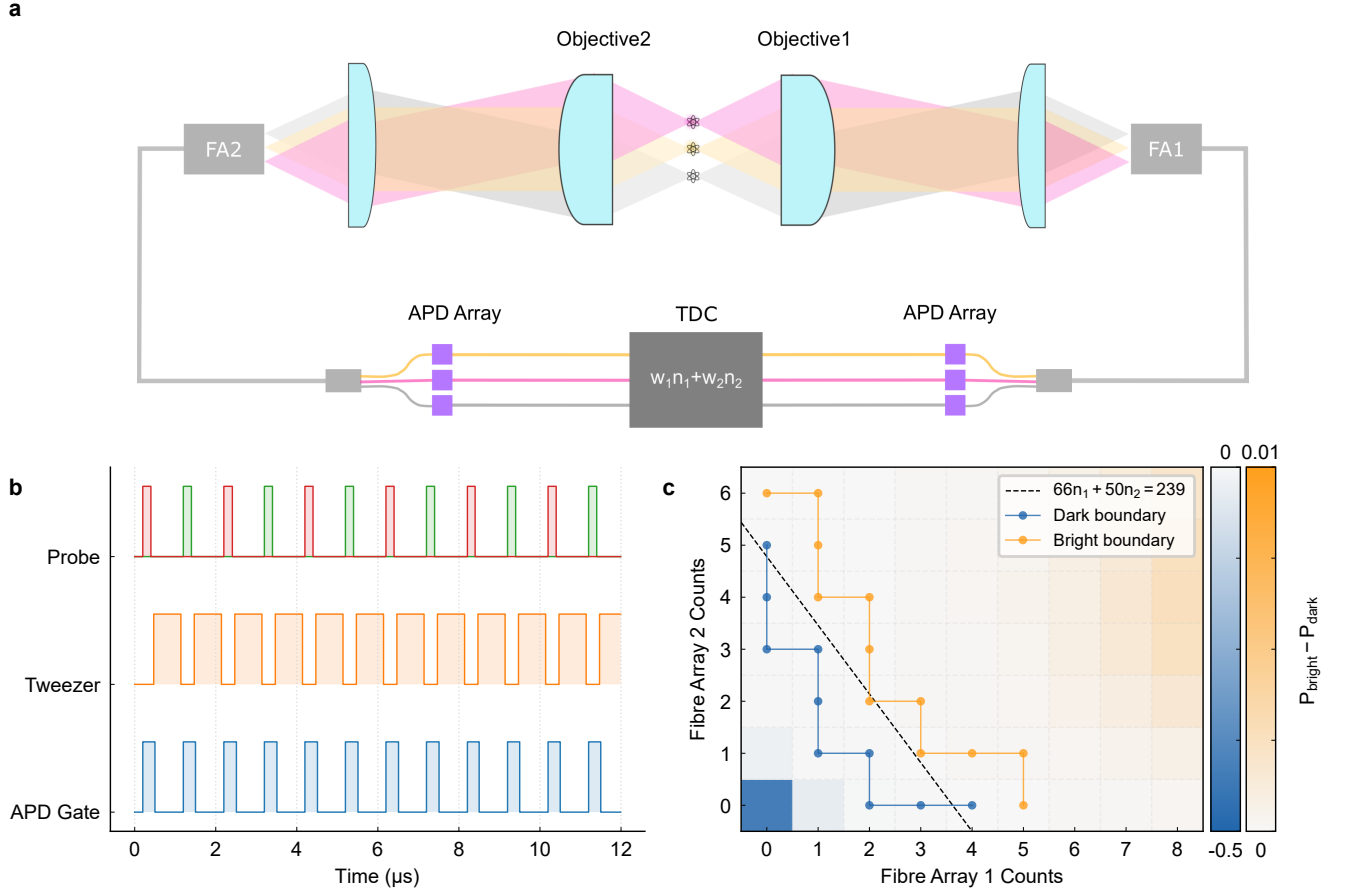


Figure S1. **Dual-FA state identification.** (a) Dual-objective setup for the fluorescence readout. Fluorescence photons from the atom array are collected by two confocal objectives, coupled into fiber arrays (FA1 and FA2), and detected by APD arrays for site-resolved single-photon detection. The two photon-count streams are combined in a time-to-digital converter (TDC) using collection-efficiency-dependent weights ($w_1 n_1 + w_2 n_2$) to optimize readout performance. (b) Detailed FND sequence. The red and green pulses denote the interleaved activation of Probe 1 and Probe 2, respectively. The orange trace shows the trapping tweezer operating at a 1 MHz rate, strobed out of phase with the pulsed probe beams. The blue trace indicates the temporal APD gate, synchronized with the probe windows to filter out background noise. (c) Joint photon-count distribution and state-discrimination boundary. The heat map displays the joint photon counts probability distribution measured from two fibre arrays. Each cell represents the difference between the bright- and dark-state probabilities, with dark-state-dominated regions shown in blue and bright-state-dominated regions shown in orange. The black dashed line denotes the empirically optimized linear decision boundary.

The atom is identified as being in the bright state when the bright-state likelihood exceeds the dark-state likelihood, $P(n_1, n_2|B) > P(n_1, n_2|D)$. Taking the logarithm of this likelihood ratio, the decision criterion becomes

$$\ln \frac{P(n_1, n_2|B)}{P(n_1, n_2|D)} > 0.$$

Substituting the Poisson likelihoods into this inequality yields

$$\sum_{i=1}^2 \left[n_i \ln \left(\frac{\mu_{B,i}}{\mu_{D,i}} \right) - (\mu_{B,i} - \mu_{D,i}) \right] > 0,$$

Rearranging this inequality naturally recovers the linear decision boundary introduced above, providing exact

analytical expressions for the weighting factors and the threshold:

$$w_i = \ln \left(\frac{\mu_{B,i}}{\mu_{D,i}} \right), \quad C = \sum_{i=1}^2 (\mu_{B,i} - \mu_{D,i}).$$

In our experimental configuration, the measured mean photon numbers for the fibre array 1 are $\mu_{B,1} \approx 14.4$ and $\mu_{D,1} \approx 0.1$, while those for the fibre array 2 are $\mu_{B,2} \approx 4.8$ and $\mu_{D,2} \approx 0.06$. Calculating the exact MLE weights using these values yields floating-point numbers. However, performing floating-point arithmetic is highly inefficient for real-time processing on a FPGA. Furthermore, the actual experimental photon count distribution is not strictly Poisson (see below). To address these prac-

tical constraints, we empirically optimized the estimator into an integer form:

$$66n_1 + 50n_2 > 239,$$

which was obtained from the measured fluorescence distributions of the two fiber arrays, as illustrated in Fig. S1(c).

Markov tree of benchmark errors. To assign a rare error string to a physical error channel, we construct a Markov tree of possible atomic state trajectories. Each node is written as

$$v_j = (r_j, s_j),$$

where $r_j \in \{0, 1\}$ is the measured outcome of the corresponding stage, with 1 denoting an above-threshold signal and 0 denoting a sub-threshold signal, and s_j is the actual atomic state at the end of that stage. For initialization and FNDR stages, $s_j \in \{B, D, L\}$, where B , D , and L denote bright, dark, and lost atoms, respectively. For the atom-presence measurement, the repump light mixes the bright and dark hyperfine manifolds, so the final state is represented only as $s_j \in \{B/D, L\}$.

The Markov tree includes two classes of error processes.

First, the atomic state can change during an FNDR stage through

$$B \rightarrow D, \quad B \rightarrow L, \quad D \rightarrow L, \quad D \rightarrow B,$$

with probabilities P_{depump} , P_{BL} , P_{DL} , and P_{repump} , respectively. During an atom-presence measurement, where B and D are not distinguished, the only state-change channel retained in the tree is

$$B/D \rightarrow L,$$

with probability P_L .

Second, a readout error can occur when the recorded binary outcome is inconsistent with the atomic state. Since the atomic state can change during the readout, which already corresponds to a state-change error discussed above, we attribute a readout error only when the readout result disagrees with the expectations from both the beginning atomic state and the final atomic state. For an FNDR stage, the expected readout is 1 for B and 0 for D or L . We therefore label a $B \rightarrow B$ edge followed by a measured 0 as a bright-assignment error P_{BE} , and a $D/L \rightarrow D/L$ edge followed by a measured 1 as a dark-assignment error P_{DE} . For an atom-presence measurement, the expected readout is 1 for B/D and 0 for L , with the same definitions of P_{BE} and P_{DE} applied to the corresponding presence/absence discrimination. If a state-change error and a readout error are both required on the same trajectory, the corresponding edge is labelled by the product of probabilities.

For a measured string $\mathcal{S}$, we retain all trajectories $\gamma \in \Gamma(\mathcal{S})$ whose node outcomes reproduce that string. The

weight of a trajectory is the product of the edge labels,

$$W(\gamma) = \prod_{e \in \gamma} w(e),$$

where $w(e) = 1$ for an error-free transition and $w(e) \in \{P_{\text{depump}}, P_{\text{BL}}, P_{\text{DL}}, P_{\text{repump}}, P_L, P_{\text{BE}}, P_{\text{DE}}\}$ for an error edge. Since all benchmarked error probabilities are small, the observed probability of a given string is assigned to the lowest non-vanishing order,

$$P(\mathcal{S}) = \sum_{\gamma \in \Gamma(\mathcal{S})} W(\gamma) \simeq \sum_{\gamma \in \Gamma(\mathcal{S}), n(\gamma)=n_{\min}} W(\gamma),$$

where $n(\gamma)$ is the number of error labels along the trajectory and $n_{\min}$ is the minimum order required to explain the measured string. Higher-order trajectories are shown in the tree, but they are not included in the first-order rate extraction.

For the bright-error benchmark, the atom is first initialized into the bright manifold by optical pumping, and then measured by three consecutive FNDR stages. The tree is conditioned on the measured string $B101$, where the leading initial branch is B with probability $O(1)$, while imperfect preparation into D or L enters as higher-order branches with probabilities P_{depump} and P_{BL} . The lowest-order trajectory consistent with $B101$ is

$$B \xrightarrow{\text{FNDR } 1} (1, B) \xrightarrow{\text{FNDR } 2} (0, B) \xrightarrow{\text{FNDR } 3} (1, B),$$

where the second FNDR returns a sub-threshold signal even though the atom remains bright. This trajectory has weight P_{BE} and is assigned to the bright-assignment error. Trajectories involving an initial depump followed by a repump, or loss followed by a dark-error event, are at least second order and are neglected in the extracted first-order rate.

Similarly, for the dark-assignment-error benchmark, the atom (verified existed at the beginning) is first initialized into the dark manifold by optical depumping, and then measured by three consecutive FNDR pulses. The tree is conditioned on the measured string $D010$. The dominant initial branch is D with probability $O(1)$, while residual preparation in B and preparation loss enter with probabilities $O(P_{\text{repump}})$ and P_{DL} . The lowest-order trajectory consistent with $D010$ is

$$D \xrightarrow{\text{FNDR } 1} (0, D) \xrightarrow{\text{FNDR } 2} (1, D) \xrightarrow{\text{FNDR } 3} (0, D),$$

where the second FNDR returns an above-threshold signal although the atom remains dark. This trajectory has weight P_{DE} and is assigned to the dark-assignment error.

For the bright-loss benchmark, the atom is first prepared in the bright manifold by optical pumping, followed by one FNDR pulse and two atom-presence measurements. The tree is conditioned on the measured string $B100$. The leading FNDR-induced loss trajectories are

$$B \xrightarrow{\text{FNDR}} (1, L) \xrightarrow{\text{presence } 1} (0, L) \xrightarrow{\text{presence } 2} (0, L),$$

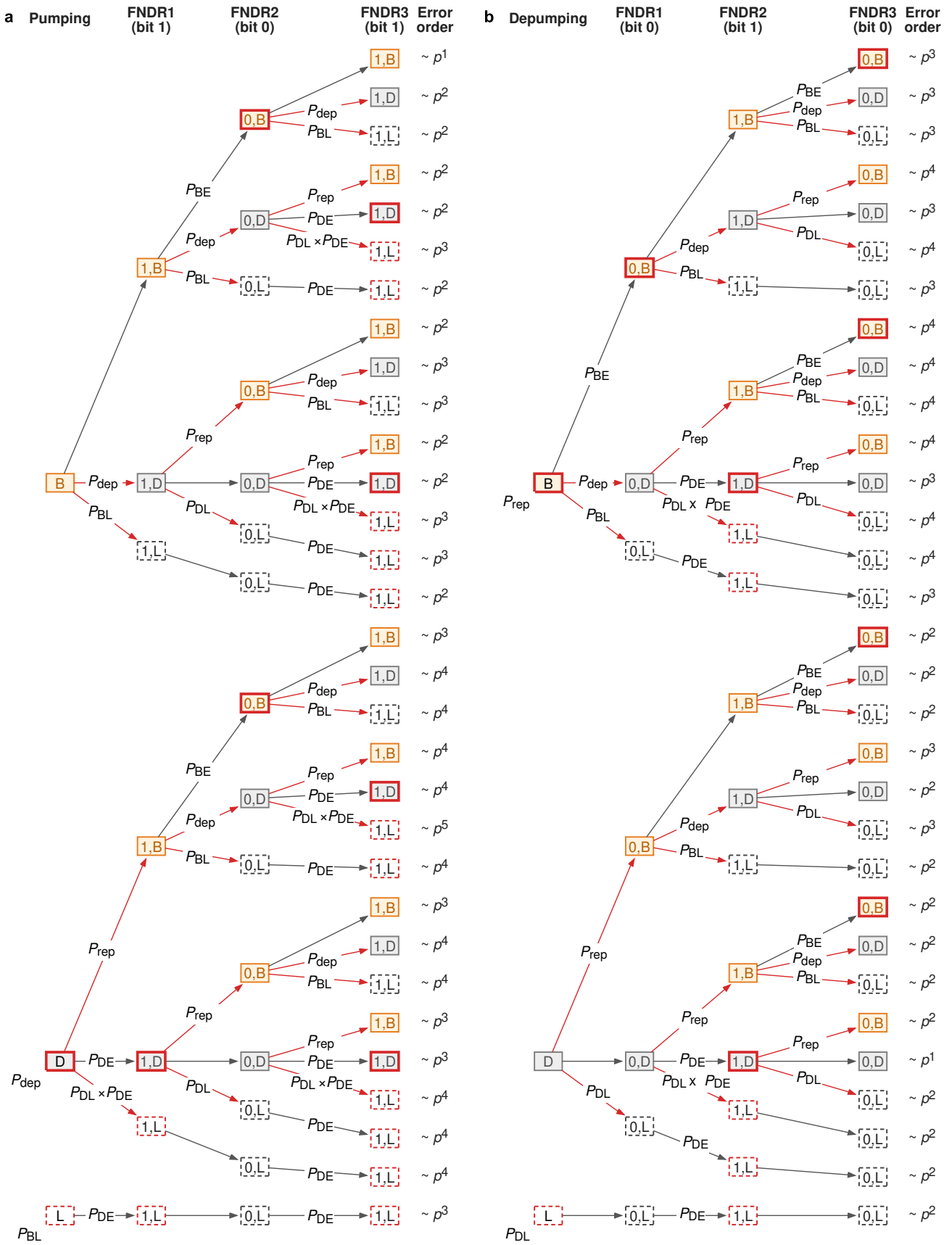


Figure S2. Markov tree for error benchmarking. (a) Bright error Markov tree. (b) Dark error Markov tree

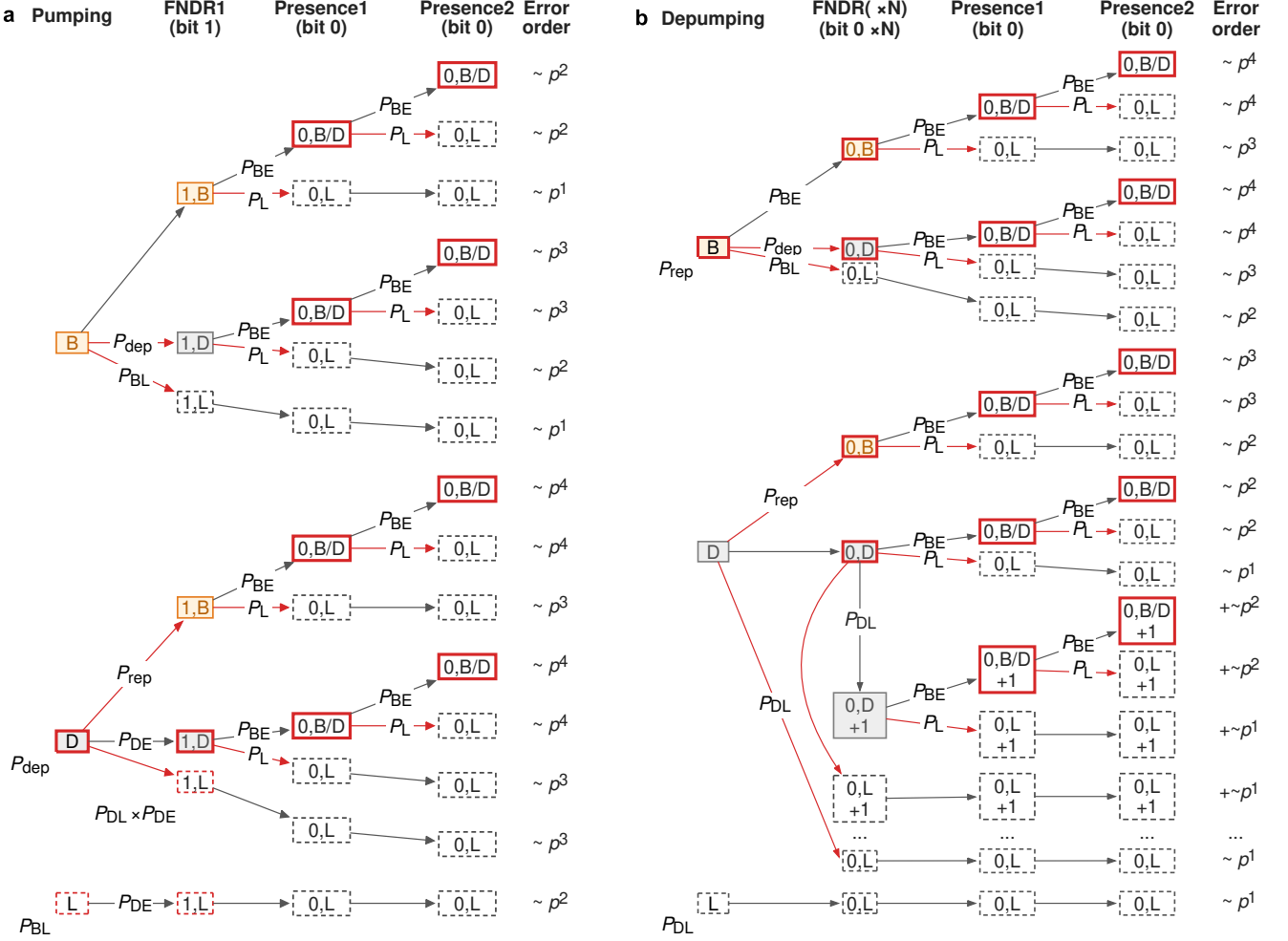


Figure S3. **Markov tree for loss benchmarking.** (a) Bright-state loss Markov tree. (b) Dark-state loss Markov tree

with weight P_{BL} . The first FND R result can still be bright because fluorescence may be collected before or during the loss event, while the two subsequent presence measurements identify the atom as absent. The Markov tree also contains a competing trajectory in which the atom survives FND R and is lost during a later presence measurement, with weight P_L . This path also contributes to our reported FND R-induced bright loss rate, and therefore the reported value represents an upper limit of the actual loss rate.

For the dark-loss benchmark, the atom (existence verified at the beginning) is first prepared in the dark manifold by optical depumping, followed by an FND R block and two atom-presence measurements. The tree is conditioned on the measured string $D0^{xN}00$. The leading FND R-induced dark-loss trajectory that increase with N is

$$D \xrightarrow{\text{FND R}} \dots \xrightarrow{\text{FND R}} (0, D) \xrightarrow{\text{FND R}} (0, L) \xrightarrow{\text{FND R}} \dots$$

$$\xrightarrow{\text{presence 1}} (0, L) \xrightarrow{\text{presence 2}} (0, L)$$

with weight P_{DL} . In the plotted tree, we keep an auxiliary label “+” to represents these trajectories that appears with increasing N (higher order trajectories are omitted for visual clarity). Similarly, the Markov tree also contains a competing trajectory in which the atom survives the increased FND R and is lost during a later presence measurement, with weight P_L . This path also contributes to our reported FND R-induced dark loss rate, and therefore the reported value represents an upper limit of the actual loss rate.

Depumping error benchmark. In this benchmark, depumping refers to probe-induced leakage from the nominally closed cycling transition, distinct from the depumping process mediated jointly by the hiding light and the probe beam discussed in Section VI. This probe-induced depumping arises from far-off-resonant excitation: residual misalignment with the quantization axis and small polarization impurity can weakly drive the $F = 2 \rightarrow F' = 2$ transition, opening leakage channels that lead to state-assignment errors when leakage occurs before sufficient photons are collected. To quantify the

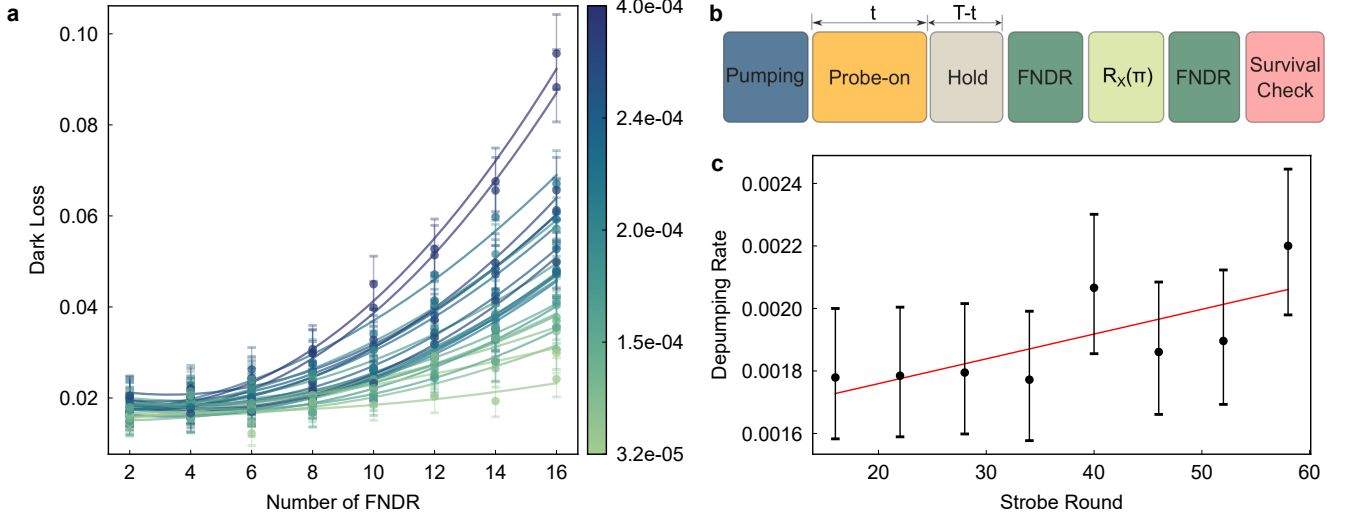


Figure S4. **Benchmarks of dark-state loss and probe-induced depumping.** (a) Dark-state loss probability for each atom, measured as a function of the number of FNDR measurements, conditioned on dark-state preparation. A quadratic fit to the averaged data yields a dark-state loss probability of $1.9^{+0.1}_{-0.1} \times 10^{-4}$. (b) Benchmark protocol for quantifying probe-induced depumping during FNDR. After pumping, a variable-duration probe-on stage is inserted before triple-state measurement, in which the two FNDR readouts are separated by a π pulse on the $|F=2, m_F=2\rangle \leftrightarrow |F=1, m_F=1\rangle$ transition. During the hold stage, the tweezer-strobing pattern is maintained while the probe light is switched off. (c) Survival-conditioned depumping probability measured as a function of the number of readout strobe rounds. A linear fit gives a depumping probability of 7.93×10^{-6} per strobe round. For the adaptive FNDR readout, an effective 37 strobe rounds gives a probe-induced depumping probability of $(2.9^{+1.0}_{-1.0}) \times 10^{-4}$ per FNDR probe stage. In (a) and (c), error bars denote 95% confidence intervals for the data points, and quoted uncertainties denote 68.3% confidence intervals from the fits.

depumping error induced by probe beams during FNDR, we insert a readout sequence after pumping and vary the duration for the probe-on stage (Fig. S4(b)). Atom loss is identified by the final survival check and treated separately as an erasure error. The following analysis therefore includes only atoms that survive the full sequence.

We denote the bright and dark readout outcomes as 1 and 0, respectively, and define N_{ij} as the number of survival-conditioned trials with outcome i in the first FNDR and outcome j in the second FNDR. The ideal case without depumping is N_{10} : the atom is bright during the first readout and is subsequently transferred to the dark state by the π pulse before the second readout. Conversely, the N_{01} and N_{00} events indicate first-order depumping error. Specifically, N_{01} corresponds to depumping to $|F=1, m_F=1\rangle$, which is reversed by the subsequent π pulse, whereas N_{00} represents leakage into other dark states that remain dark after π pulse. To ensure conservative estimate, we attribute all N_{00} events to depumping, even though some may originate from the misidentification of an initial bright state. We therefore estimate the survival-conditioned depumping probability as

$$P_{\text{depump}|\text{survive}} \leq \frac{N_{01} + N_{00}}{N_{01} + N_{00} + N_{10}}.$$

The inequality also accounts for the exclusion of the N_{11}

from the normalization. This branch is not uniquely attributable to depumping: it can arise from atoms occupying bright states outside the addressed $|F=2, m_F=2\rangle$ manifold, for example because of imperfect polarization purity, which remain bright during the first readout but are not transferred by the π pulse. Excluding this ambiguous branch reduces the denominator and therefore yields a conservative upper bound. Then we obtain 2.9×10^{-4} for the depumping probability per FNDR probe stage (Fig. S4(c)).

Background photon count. The experimental background consists of detector dark counts and stray beam scattering. The mean joint background counts is approximately 0.1 counts per readout. While the APD dark count rate is low (~ 65 cps), the afterpulsing effect introduces a non-Poissonian bias to photons distribution[5]. Specifically, this effect arises when charge carriers from a primary avalanche are trapped by lattice defects and subsequently released, triggering correlated delayed avalanches that artificially inflate the photon count.

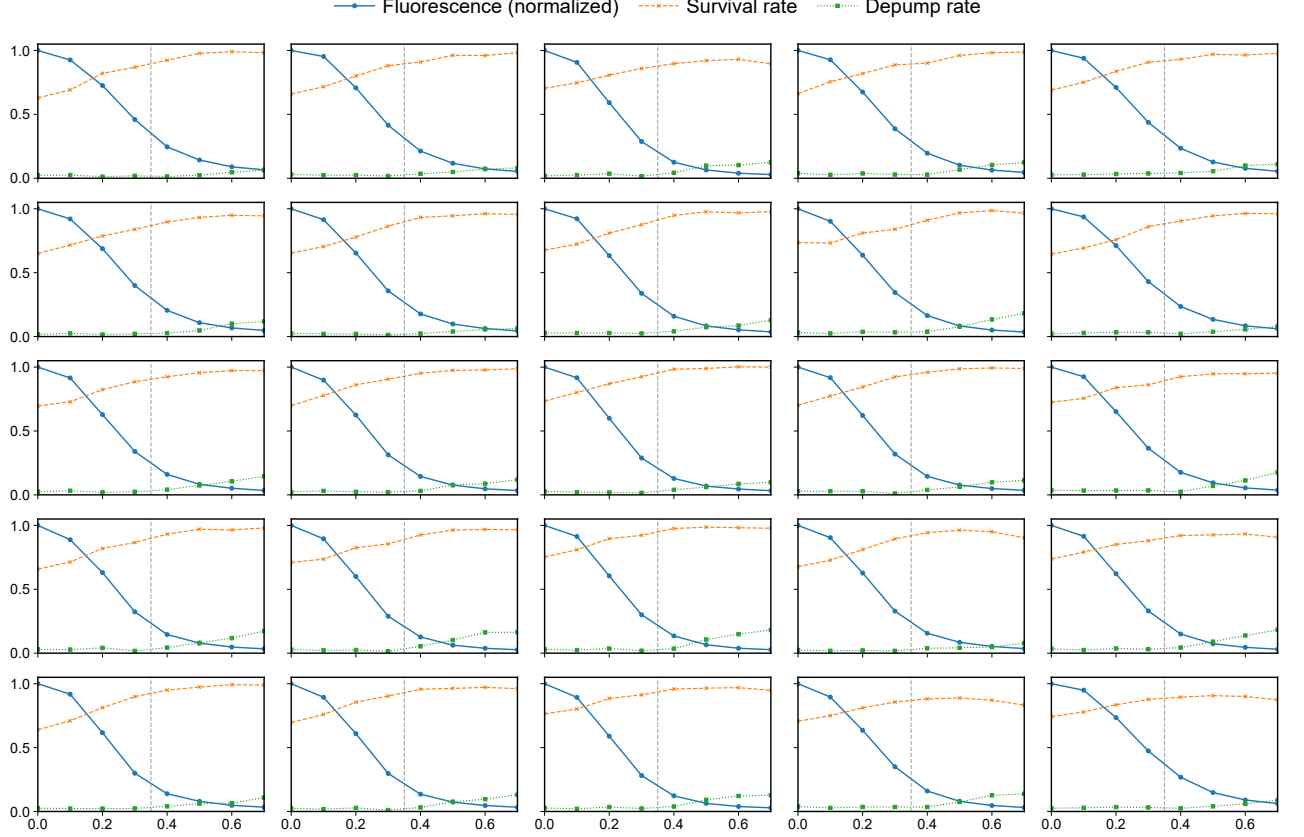


Figure S5. **hiding-beam intensity optimization.** Normalized fluorescence, depumping rate and survival rate measured for 25 individually addressed atoms after approximately 10 readout cycles. The operating amplitude of 0.5, which corresponding to approximately 100 MHz ac Stark shift, is chosen to balance atom survival and depumping.

III. MM-FA AND MLA

Fibre array design. The detection system employs two customized fibre arrays (FAs), each comprising 150 multi-mode signal channels and 8 auxiliary channels. The signal fibres have an NA of 0.12, a core diameter of $50\ \mu\text{m}$, and a pitch of $127\ \mu\text{m}$, corresponding to a waist-to-spacing ratio of approximately 1:5. The fibre pitch is positioned with a positional tolerance of $\pm 1\ \mu\text{m}$. To enable precise co-alignment with the tweezer array and maximize collection efficiency, the eight auxiliary fibres (four single-mode and four multi-mode) are positioned at the corners of each FA, serving as global fiducial markers.

Microlens array design. The 1:5 beam-waist-to-spacing ratio required for matching with the tweezer array cannot be achieved directly with a bare polarization-maintaining fibre array (PMFA) due to its small mode diameter and large channel spacing ($250\ \mu\text{m}$ pitch). We solve this problem by integrating a microlens array (MLA) into the output facet of the PMFA. The MLA expands each output mode to a waist of $50\ \mu\text{m}$, while

preserving the desired spacing as the fibre-array pitch. The resulting Gaussian beam array is subsequently relayed through $4f$ optical systems, which demagnify the MLA plane onto the atomic plane and produce $1\ \mu\text{m}$ -waist hiding beams at the target sites. The spot size is a balanced choice between suppressing inter-site crosstalk (see below) and tolerating tiny positional offset caused by atomic motion and fabrication non-uniformity.

Intensity calibration. To determine the optimal operating amplitude of the hiding beams, we characterized the fluorescence suppression, atom survival, and depumping rates across 25 addressed sites (Fig. S5). A normalized amplitude setpoint of 0.5—corresponding to an ac Stark shift of approximately 100 MHz—yields sufficient fluorescence suppression for low-recoil readout while maintaining a tolerable depumping rate. Since this depumping can be effectively mitigated (detailed in Section VI), future implementations can employ higher beam amplitudes to achieve even greater atom survival.

Hiding beam crosstalk. We quantify the optical crosstalk of the hiding beams through the induced fluorescence suppression and the increased depumping rate.

Shifting the hiding beam by one site spacing ($5\ \mu\text{m}$) relative to the tweezer array results in negligible crosstalk.

IV. RABI CIRCUIT

State reinitialization. Prior to FNDR, the tweezer trap depth is adiabatically ramped up from 1 mK to 6 mK. After each round of FNDR, we reinitialize the atom state back to the $5S_{1/2}\ |F=2, m_F=2\rangle$ state without ramping down trap depth. This is achieved by adopting a stroboscopic optical pumping protocol that shares the timing sequence and the probe frequency of the FNDR, accompanied by a simultaneous repump beam tuned to the free-space $|F=1\rangle \rightarrow |F'=2\rangle$ transition. This resolves three operational challenges. First, continuous optical pumping at 6 mK is unfeasible, as the pronounced tweezer-induced light shifts exceed the frequency tuning range of our repump light. Second, performing the reinitialization at 1 mK trap depth is undesirable, as it would introduce a millisecond-scale ramping overhead before readout. Finally, matching the pumping light's frequency to that of the readout eliminates the $\sim 100\mu\text{s}$ delay associated with switching between different laser frequencies.

Stage	Operation	Time (μs)	
Recooling	settling time	130.0	930.0
	grey molasses	800.0	
Ramp Tweezer	settling time	110.0	910.0
	ramp up	400.0	
	ramp down	400.0	
Pumping	settling time	25.0	60.2
	probe	7.0	
	non-overlap delay	2.6	
	strobing hold	25.6	
FNDR	settling time	25.0	153.0
	probe	25.6	
	non-overlap delay	9.3	
	strobing hold	93.1	

Table SI. **Time budget for the repeated Rabi circuit.** Durations of the constituent operations of the recooling, tweezer-ramping, optical-pumping and FNDR stages. The total cycle duration is obtained by summing the durations of all operations, with the FNDR stage counted twice for triple-state measurement, and corresponds to an experimental repetition rate of 1.7 kHz. Settling times arise from hardware latencies, whereas non-overlap delays temporally separate the probe and tweezer pulses, preventing simultaneous illumination that would otherwise induce additional depumping.

Time budget. In this work, one Rabi circuit cycle has an effective duration of approximately $600\ \mu\text{s}$ (Fig. S6), with the detailed operation times shown in Table SI. This time budget is constrained by three major overheads. First, recooling introduces both a trap-ramping penalty

and a non-optimized cooling duration. At a trap depth of 6 mK, the available repump laser is far detuned from the $|F=1\rangle \rightarrow |F'=2\rangle$ transition, requiring the tweezers to be ramped down to 1 mK and thereby adding a $910\ \mu\text{s}$ overhead. Additionally, the active cooling sequence currently lasts $800\ \mu\text{s}$, although $500\ \mu\text{s}$ is experimentally sufficient. Second, latencies associated with laser-frequency switching and mechanical-shutter actuation accumulate as preparatory settling times throughout the circuit sequence. Third, the tweezers remain on for 67% of the FNDR and optical-pumping intervals, further contributing to the overall time cost.

Future implementations could systematically address these constraints. Introducing a repump laser resonant with the $|F=1\rangle \rightarrow |F'=2\rangle$ transition at 6 mK would remove the $910\ \mu\text{s}$ ramping overhead and enable efficient optical pumping at this elevated depth, reducing the pumping stage to the probe-limited duration. Furthermore, dedicated laser beams pre-tuned to the readout and pumping frequencies, together with independent shutters, would eliminate settling times. Optimizing the tweezer-to-probe duration towards a 1:1 ratio will further shorten the FNDR blocks. Together, these improvements are expected to increase the circuit clock rate to approximately 5 kHz.

V. GREY MOLASSES COOLING

Grey molasses cooling during loading. After the MOT-loading stage, we perform 1D grey-molasses cooling for 6 ms. The π -polarized cooling beam is collinearly combined with a mixed-polarization repump beam. The cooling beam is blue-detuned by 94.2 MHz from the $D_2\ |F=2, m_F=2\rangle \rightarrow |F'=2, m'_F=2\rangle$ transition. To reduce spatial inhomogeneity in the cooling, we impose a 5 kHz relative frequency offset between the two counter-propagating cooling beams, thereby translating the standing-wave intensity pattern during the cooling interval. With this configuration, the loading-stage grey molasses cools the atoms to approximately $14\ \mu\text{K}$, as measured by release-and-recapture thermometry.

Grey molasses cooling. During the recooling stage, we perform grey-molasses cooling in the 10 G bias magnetic field that defines the quantization axis. We retain the loading-stage grey-molasses beam arrangement and add a pair of counter-propagating σ^+ -polarized cooling beams along this axis. The cooling beams are blue-detuned by 54.3 MHz from the $D_2\ |F=2, m_F=2\rangle \rightarrow |F'=2, m'_F=2\rangle$ transition, with the value chosen by experimental optimization. Under these conditions, the recooling step reduces the atomic temperature to approximately $20\ \mu\text{K}$ in $800\ \mu\text{s}$.

To interpret the cooling observed at finite magnetic field, we develop a semiclassical dressed-state model. The bias magnetic field lifts the degeneracy of the m_F sub-

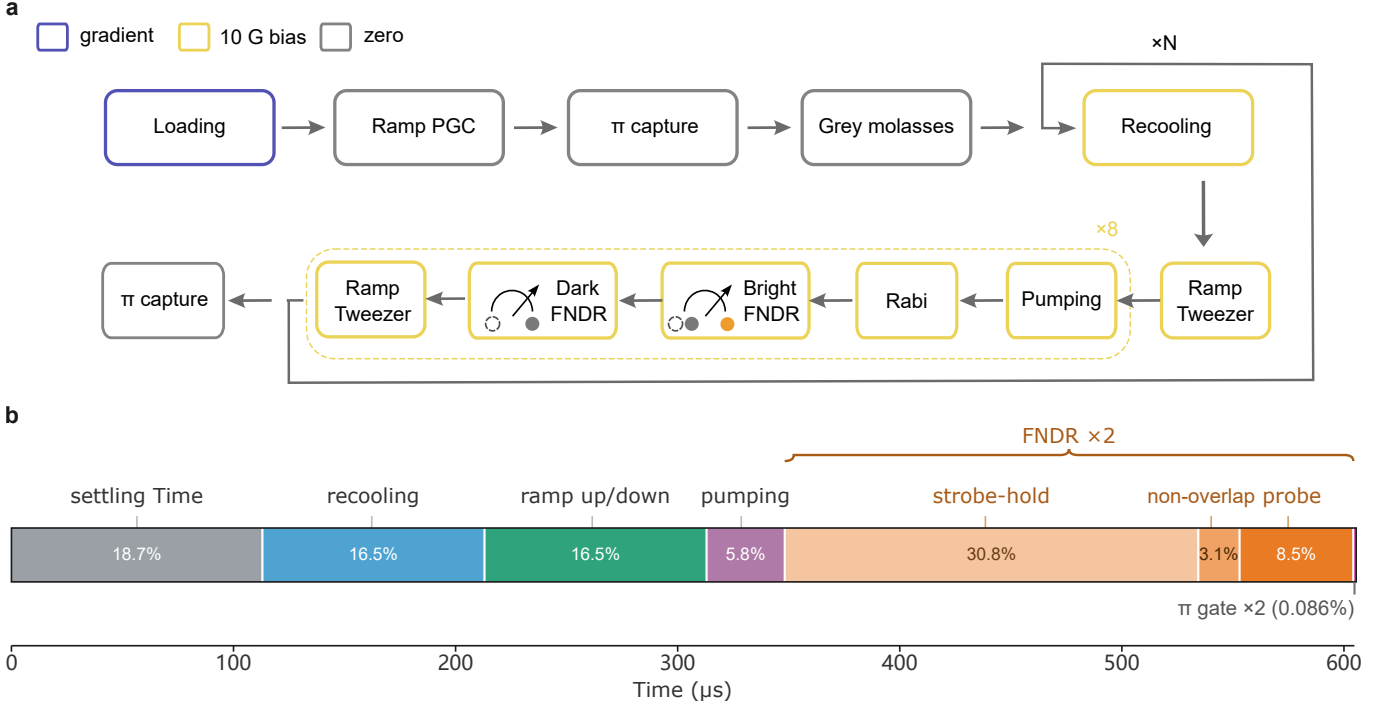


Figure S6. **Experimental sequence and time budget analysis of rabi circuit.** (a) Experimental sequence for atom preparation, recooling, repeated Rabi circuit, and FNDR measurements. The box colors denote the magnetic-field configuration: blue, magnetic-field gradient; grey, zero magnetic field; yellow, 10G bias field. (b) Detailed time budget breakdown for a single Rabi circuit cycle operating at 1.7 kHz repetition rate. The probe time in FNDR is highly compressed, accounting for merely 8.5% of the total duration. By further optimizing experimental parameters (detailed in Section V), the repetition rate can potentially be increased to 5 kHz.

levels, so the standard zero-field grey-molasses picture is no longer directly applicable [6, 7]. In this regime, atoms couple to moving intensity lattices generated by the π - and σ^+ -polarized beam pairs. The atom-light couplings produce ground-like dressed-state channels whose light shifts and scattering rates vary with position. During the cooling cycle, atoms climb a low-scattering dressed channel and are optically pumped near high-scattering regions into a lower-energy channel, thereby losing kinetic energy through a Sisyphus-type grey-molasses mechanism.

We support this interpretation with a semiclassical Monte Carlo simulation of the coupled internal-state and motional dynamics, as summarized in Fig. S7. The model includes the five Zeeman sublevels of the $F = 2$ ground manifold and the five Zeeman sublevels of the $F' = 2$ excited manifold. At each position and time, the single-atom Hamiltonian is

$$H(\mathbf{r}, t) = H_{\text{diag}} + E_{\pi}(x, t)H_{\pi} + E_{\sigma^+}(z, t)H_{\sigma^+} + H_{\text{tw}}(\mathbf{r}),$$

where H_{diag} contains the Zeeman energies and optical detuning, and H_{tw} denotes the tweezer-induced light shift. The matrices H_{π} and H_{σ^+} contain the Clebsch-Gordan-weighted dipole couplings for the two polarizations, while $E_{\pi}(x, t)$ and $E_{\sigma^+}(z, t)$ describe the spatiotemporal am-

plitudes of the moving lattices. The density matrix obeys

$$\dot{\rho} = -\frac{i}{\hbar}[H, \rho] + \sum_{\ell} \left(C_{\ell} \rho C_{\ell}^{\dagger} - \frac{1}{2} \{ C_{\ell}^{\dagger} C_{\ell}, \rho \} \right),$$

where the collapse operators C_{ℓ} describe spontaneous emission between the excited- and ground-state Zeeman sublevels. The external motion is propagated semiclassically using the Ehrenfest force

$$m\ddot{\mathbf{r}}_i = -\text{Tr}[\rho \partial_i H(\mathbf{r}, t)],$$

with stochastic recoil kicks added according to the instantaneous excited-state population. Figure S7 (a) shows the spatial structure underlying the cooling mechanism. The calculated dynamics show population redistribution among the ground-like dressed channels, together with a decrease in kinetic energy during the cooling interval (Fig. S7 (b)). A simulated detuning scan gives optimal cooling near 50 MHz, consistent with the experimentally optimized detuning of 54.3 MHz. The final temperature predicted by the simulation is lower than the measured value, which we attribute to the idealized two-dimensional model and to omitted heating mechanisms, including axial scattering and recoil heating, residual polarization imperfections, magnetic-field inhomogeneity, and other technical noise.

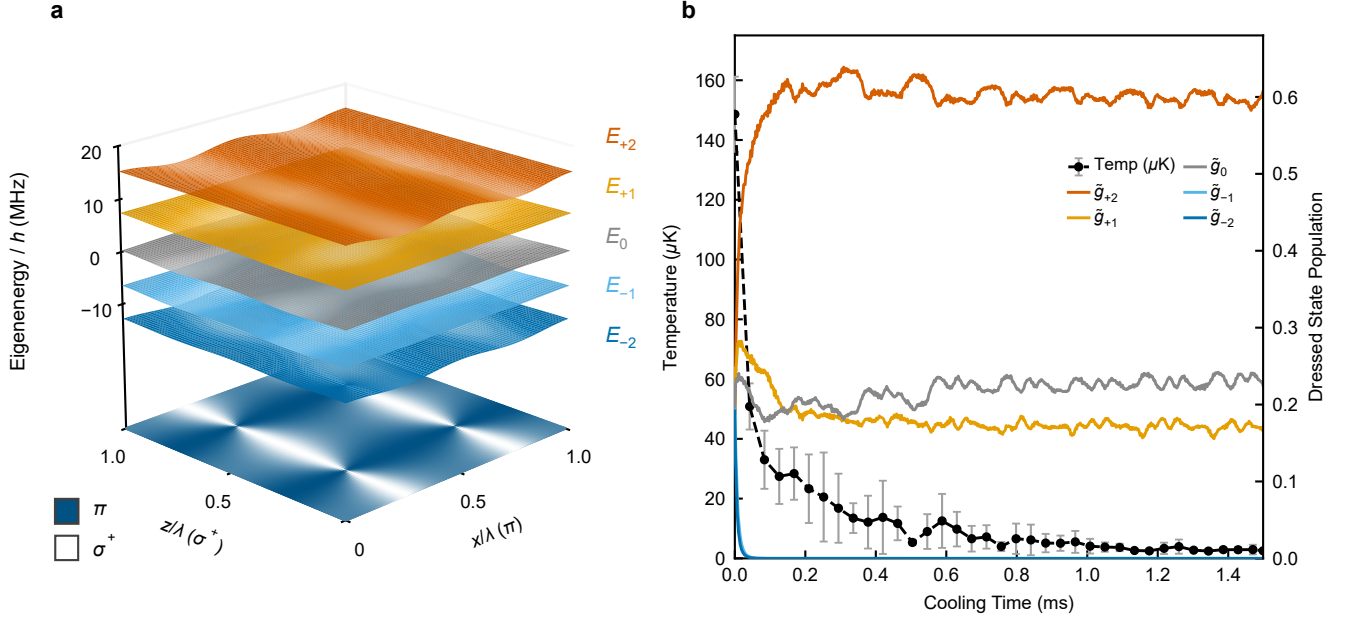


Figure S7. **Numerical simulation of finite-field grey-molasses cooling.** (a) Ground-like dressed-state energy surfaces and the spatial pattern of the local π - and σ^+ -polarized components for a representative lattice phase. Blue and white denote the π and σ^+ components, respectively. The branches E_{-2} – E_{+2} define the adiabatic channels involved in the Sisyphus-type grey-molasses mechanism. (b) Population projected onto the ground-like dressed channels $|\tilde{g}_{m_F}\rangle$ (right axis) and ensemble-averaged kinetic temperature (left axis) as functions of cooling time. Each label $|\tilde{g}_{m_F}\rangle$ denotes the dressed state with dominant overlap with the corresponding bare state $|F=2, m_F\rangle$, for $m_F = -2, -1, 0, 1, 2$. The temperature is evaluated at the experimentally optimized detuning $\Delta/2\pi = 54.3$ MHz. In the idealized simulation, it reaches 3.0 μK after 1.5 ms.

VI. SIMULATION FOR HIDING BEAM

We use a local hiding beam to implement site-resolved adaptive protection, thereby suppressing redundant scattering. Hiding-beam-induced state transfer has different consequences depending on when and where it occurs. At a site classified as bright, the hiding beam is activated only after the measurement outcome has been determined and therefore cannot affect the state assignment in that measurement. At sites that have not yet been classified, however, residual hiding light arising from finite on/off extinction and inter-site crosstalk can induce depumping before the decision is reached, and therefore causes readout error.

852-nm hiding implementation. For the present 852-nm implementation, we consider three depumping mechanisms.

The first is spontaneous Raman scattering (SRS) driven by the far-off-resonant, nominally π -polarized hiding beam. Destructive interference between the D1- and D2-mediated scattering amplitudes suppresses this contribution, with a rate proportional to $\left|\frac{1}{\Delta_{D1}} - \frac{1}{\Delta_{D2}}\right|^2$ [8]. At the on-site hiding beam intensity used in this experiment, the corresponding depumping rate is estimated to be approximately 1.5 s^{-1} .

The second mechanism results from probe-light polar-

ization impurity. The differential ac Stark shift produced by the 852-nm hiding beam reduces the detuning of the 780-nm probe from the $|5S_{1/2}, F=2\rangle \leftrightarrow |5P_{3/2}, F'=2\rangle$ transition. For an atom initially in $|F=2, m_F=2\rangle$, this transition is forbidden under ideal σ^+ -polarized probe illumination. Residual π and σ^- components of the probe, however, can weakly drive off-resonant excitation through the $F'=2$ manifold, followed by spontaneous decay into the dark hyperfine manifold. We refer to this process as off-resonant imaging-induced scattering (OIS).

The third mechanism is Raman-assisted depumping (RAD). Tight focusing and residual alignment errors give the local hiding field finite π , σ^+ , and σ^- components relative to the quantization axis. These components open V-type Raman-assisted depumping pathways from the cycling excited state $|5P_{3/2}, F'=3, m_{F'}=3\rangle$ to states outside the closed transition. As the hiding-beam power increases, differential ac Stark shifts bring the relevant excited-state levels close to Raman resonance, giving rise to the sharp depumping feature shown in Fig. S8 (b).

To quantify these mechanisms, we employ a master-equation model that includes the relevant ground- and excited-state hyperfine and Zeeman sublevels. The model incorporates the measured probe-beam scattering rate, the calibrated hiding-beam-induced detuning, the estimated polarization impurity, and the residual hiding

beam illumination arising from the finite extinction ratio and crosstalk. The modeling is detailed in a later subsection. The calculated contributions from the three mechanisms described above are summarized in Fig. S8 (b).

1529.6-nm hiding implementation. A prospective alternative is to use 1529.6-nm hiding light red-detuned from the $5P_{3/2} \rightarrow 4D$ manifold [9–12]. Compared with the present 852-nm implementation, this scheme offers three advantages. It produces a large light shift of the $5P_{3/2}$ manifold while leaving the $5S_{1/2}$ ground manifold nearly unaffected, preserves the $F' = 3 - F' = 2$ excited-state separation, and requires substantially less optical power. The remaining depumping channels are Raman-assisted transfer and two-photon excitation through the $4D$ manifold. The latter can be reduced by increasing the detuning from the $5P_{3/2} \rightarrow 4D$ transition.

The numerical results are shown in Fig. S8 (d). At the on-site operating power required to suppress probe scattering by a factor of ten, hiding-beam-induced depumping occurs only after the site has been classified and therefore does not affect the state assignment. The relevant error constraint instead applies to the residual hiding light incident on sites that have not yet been classified. Requiring the corresponding depumping rate to remain below $\sim 1 \text{ s}^{-1}$ limits the resulting depumping probability during the pre-decision readout interval to below 10^{-4} . This condition requires the combined residual intensity arising from inter-site crosstalk and finite on/off extinction to be at least 35 dB below the on-site hiding intensity, which is compatible with current experimental techniques.

Master-equation model. To quantitatively evaluate the depumping mechanisms, we use an open-system density-matrix model that includes the relevant hyperfine and Zeeman sublevels of the $5S_{1/2}$, $5P_{3/2}$, and, for the 1529.6-nm scheme, $4D$ manifolds. Each basis state is denoted by

$$|\alpha\rangle = |nL_J, F, m_F\rangle.$$

For the 852-nm calculation, we include the relevant hyperfine and Zeeman sublevels of the $5S_{1/2}$ ground manifold and the $5P_{3/2}$, $F' = 2, 3$ excited manifolds. For the 1529.6-nm calculation, the basis is extended to include the relevant sublevels of the $5S_{1/2}$, $5P_{3/2}$, $4D_{3/2}$, and $4D_{5/2}$ manifolds.

For a given local hiding-beam power P , the density matrix is evolved according to the Lindblad master equation

$$\dot{\rho} = -\frac{i}{\hbar}[H, \rho] + \sum_{\mu} \left(L_{\mu} \rho L_{\mu}^{\dagger} - \frac{1}{2} \{ L_{\mu}^{\dagger} L_{\mu}, \rho \} \right).$$

In the rotating-wave approximation, the Hamiltonian is written as

$$\frac{H}{\hbar} = \sum_{\alpha} \Delta_{\alpha}(P) |\alpha\rangle \langle \alpha| + \frac{1}{2} \sum_{\alpha, \beta} [\Omega_{\alpha\beta} |\alpha\rangle \langle \beta| + \text{H.c.}]$$

Here, Δ_{α} includes the bare probe detuning, the hyperfine and Zeeman splittings, and the hiding-beam-induced Stark shift. The latter is obtained from the calibrated differential light shift; equivalently, its power dependence is calculated from the dynamic polarizability as

$$\delta_{\alpha}^{\text{ac}}(P) = -\frac{\alpha_{\alpha}(\omega_h) I_h}{2c\epsilon_0}, \quad I_h = \frac{2P}{\pi w_0^2}.$$

For a hyperfine state, the electronic-state polarizabilities are weighted by the corresponding Clebsch–Gordan coefficients,

$$\alpha_{Fm_F} = \sum_{m_J, m_I} |\langle Jm_J, Im_I | Fm_F \rangle|^2 \alpha_{Jm_J}.$$

The optical coupling between two sublevels is calculated from the dipole matrix element,

$$\Omega_{\alpha\beta}^{(q)} = \frac{E_q}{\hbar} \langle \beta | er_q | \alpha \rangle, \quad E_q = \sqrt{\frac{4P_q}{\pi w_0^2 c \epsilon_0}},$$

where $q = 0, \pm 1$ denotes the π , $\sigma^{\pm}$ polarization component. The relative coupling strengths therefore include the appropriate Clebsch–Gordan and Wigner-6j factors. A polarization impurity specified by ξ_q dB is included as a power fraction

$$\eta_q = 10^{\xi_q/10}, \quad P_q = \eta_q P,$$

so that the corresponding Rabi frequency scales as $\Omega_q \propto \sqrt{\eta_q P}$. In the calculation, the total polarization impurities of the hiding and probe light are set to -20 and -40 dB, respectively.

Off-resonant hiding-beam couplings also generate effective Raman coupling between states within the $5P_{3/2}$ manifold. After adiabatically eliminating the intermediate state, the effective two-photon coupling is

$$\Omega_{ij}^{(\text{R})} = \sum_k \frac{\Omega_{ik}^{(\text{h})} \Omega_{kj}^{(\text{h})}}{2\Delta_k},$$

where k labels the intermediate sublevels and Δ_k is the corresponding one-photon detuning. As the hiding beam brings the $F' = 3$ manifold close to resonance with the neighboring $F' = 2$ manifold, these effective couplings strongly enhance Raman-assisted depumping.

Spontaneous emission is represented by collapse operators

$$L_{g \leftarrow e} = \sqrt{\Gamma_{e \rightarrow g}} |g\rangle \langle e|,$$

with hyperfine- and polarization-resolved branching rates

$$\Gamma_{e \rightarrow g} = \Gamma_{J_e \rightarrow J_g} (2F_g + 1)(2J_e + 1) \left\{ \begin{matrix} J_e & J_g & 1 \\ F_g & F_e & I \end{matrix} \right\}^2 \times |\langle F_g m_g; 1q | F_e m_e \rangle|^2.$$

Thus, the model accounts for both population redistribution among the hyperfine and Zeeman sublevels and depumping out of the protected cycling manifold.

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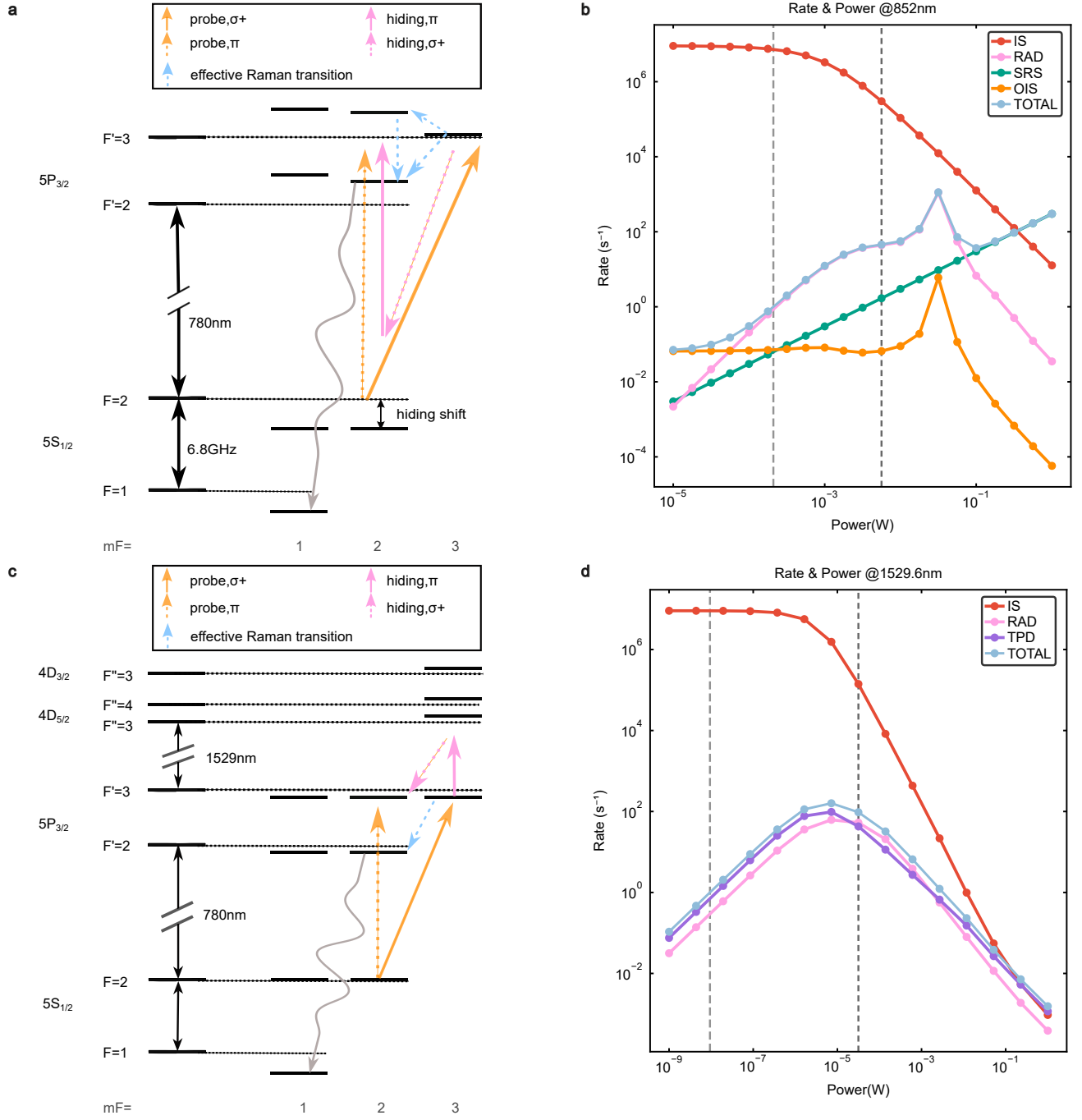


Figure S8. **Depumping mechanisms.** (a) Level structure for the 852-nm hiding scheme. Dashed and solid lines denote unshifted and light-shifted levels, respectively. The differential ac Stark shift reduces the $F' = 3 - F' = 2$ splitting and enhances polarization-impurity-induced V-type Raman coupling. (b) Calculated rates versus 852-nm hiding-beam power P . The imaging-light scattering rate (IS, red) is suppressed as P^{-2} . Hiding-beam-induced spontaneous Raman scattering (SRS, green) induced depumping increases approximately linearly with P . V-type Raman-assisted depumping (RAD, pink) proceeds through coherent transfer from $|F' = 3, m_F' = 3\rangle$ to $|F' = 2, m_F' = 2\rangle$, followed by decay into $F = 1$. Off-resonant imaging-induced scattering (OIS) through the $F' = 2$ manifold caused by residual polarization impurity is shown in orange. The total depumping rate is shown in light blue (TOTAL). The left dashed line marks a residual-power threshold with a total depumping rate of 1 s^{-1} ; the right dashed line marks a typical hiding-beam-on operating point after state discrimination. For simulation we use $w_0 = 1.0 \text{ }\mu\text{m}$, and the hiding- and imaging-light polarization impurities are -20 and -40 dB, respectively. (c) Level structure for the 1529.6-nm hiding scheme. Coupling to the $4D$ states shifts the $5P_{3/2}$ excited states while approximately preserving the $F' = 3 - F' = 2$ splitting. (d) Calculated rates versus 1529.6-nm hiding-beam power P . The RAD rate (pink) peaks near $10 \text{ }\mu\text{W}$ but remains below that of the 852-nm scheme. The two-photon depumping (TPD) rate (shown in purple) results from two-photon excitation to $4D$ followed by decay into the $F = 1$ manifold. It can be reduced by increasing the $5P_{3/2} \rightarrow 4D$ detuning. The total depumping rate is shown in light blue (TOTAL). The dashed lines have the same meanings as in (b). For simulation we use $w_0 = 2 \text{ }\mu\text{m}$, and the hiding- and imaging-light polarization impurities are -20 and -40 dB, respectively.