High-Fidelity State Detection of Alkali-Metal Atoms in Optical Tweezer Traps

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We demonstrate discrimination of ground-state hyperfine manifolds of a cesium atom in an optical tweezer using a simple probe beam with 0.16 \(\pm\) 0.02% detection infidelity and 2.6(2)% detection-driven loss of the bright state, nearly an order of magnitude better than previously published low-loss readout results for alkali-metal atoms in optical tweezers. By accepting 14.1(3)% loss of the bright state, we reach a state discrimination fidelity of 99.89(3)%. Our low-atom-loss and high-fidelity state detection takes advantage of an adaptive detection method to mitigate heating during detection and eliminates the extra depumping mechanism due to population transfer between excited-state sublevels through V-type stimulated Raman transitions caused by the trap laser when the probe laser is present. In this work, complex optical systems and stringent vacuum are not required.

Neutral atoms in arrays of optical dipole traps (ODTs) have been demonstrated to provide a promising platform for quantum computing [13], quantum simulation [4, 5], quantum chemistry [6], and optical clocks [7]. Alkali-metal or other Hydrogen-like atoms are frequently utilized because of the simplicity of their electronic structure and well-established methods of laser cooling [12].

Performance of alkali-metal atoms in tweezer platforms has to date been limited by the state readout fidelity. Various detection schemes have been used, primarily falling under state-dependent fluorescence collection or loss-based schemes where the state of the atom is mapped to trap occupation. Although relatively high detection fidelity has been achieved with loss-based schemes [3], these methods impose a vacuum-dependent upper bound on readout fidelity, slow the repetition rate, and complicate algorithms requiring mid-circuit measurement. Fluorescence detection on the other hand allows – in principle – for high fidelity state measurement without losing the atom. Previously, results for alkali atoms with fluorescence detection have only shown >1.2% infidelity without enhancement from an optical cavity [13–18]. Alternately, high-fidelity, low-loss detection of alkali atoms has been demonstrated in an optical lattice using a state-dependent potential method [19], however, adaptation of this method to optical tweezers is not straightforward. Due to both optical depumping induced by the probe and trap lasers and detection heating driven atom loss from the trapping potential, higher-fidelity state detection of alkali-metal atoms in optical tweezers using a simple probing scheme had not been demonstrated until now.

In this paper, we report state discrimination between ground-state hyperfine manifolds/multiplets of Cs atoms with 0.16 \(\pm\) 0.02% infidelity while suffering only 2.6(2)% detection-driven losses of the bright state, despite relatively low collection efficiency. Further, by accepting moderate detection-driven losses of 14.1(3)%\%, we reach a state discrimination fidelity of 99.89(3)%. We achieve these results by implementing an adaptive detection scheme to manage detection heating and mitigating a previously unreported state information loss channel stemming from simultaneous probe and trap illumination. We do not impose any stringent vacuum or optical system requirements to achieve this result, enabling straightforward integration of our techniques on contemporary alkali-metal atom tweezer platforms.

This work is an important step towards scalable, high-performance alkali-tweezer machines. While the measurement error on noisy intermediate-scale quantum (NISQ) machines is often overlooked as it does not typically scale with algorithm length, it is important to consider the scaling of the measurement error with system size. For bit string read-out, the measurement fidelity scales as \(\mathcal{F}(N)\), where \(\mathcal{F}\) is the single-atom measurement fidelity and \(N\) is the number of bits. This indicates that the order of magnitude improvement on low-loss, single-atom fidelity reported here will be crucial for larger systems. In addition to quantum computing applications, by demonstrating simultaneous high-fidelity and low loss, we improve the outlook for quantum sensing applications where the sensitivity and/or bandwidth are directly tied to atom retention capability via repetition rate.

Our experimental sequence is triggered upon loading a Cs atom from a magneto-optical trap into a tweezer, after which we immediately cool it with polarization-gradient cooling to \(\approx 2\ \mu K\). The atom is then prepared in either \(|F = 4\rangle\) ("bright") or \(|F = 3\rangle\) ("dark") hyperfine ground manifolds and the state of the atom is then read out using a near-resonant probe laser. Finally, we check for atom presence with the cooling beams.

The tweezer trap is formed by focusing up to 10 mW of 937 nm laser light through a 0.45 NA microscope objective to a spot with a 1.6 \(\mu m\) /\(e^2\) waist radius. The trap is small enough that a light-assisted collisional-blockade mechanism ensures loading no more than one atom at a time [20]. The trap wavelength of \(\lambda_{\text{trap}} = 937\ \text{nm}\) is chosen to be red-detuned from both the \(D_1\) and \(D_2\) line transitions and to have approximately the same AC Stark shift on both states (6\(S_{1/2}\) and 6\(P_{3/2}\)) of the cooling transition (i.e., "magic" wavelength for the cooling transition). The AC Stark shift from the linearly
polaredized 937 nm light on the ground state forms the trapping potential proportional to the intensity, $\Gamma$, according to \[21\]:

$$U_{\text{ACS}} \approx \frac{\pi c^2 \Gamma}{2 \omega_0} \left( \frac{2}{\Delta D_2} + \frac{1}{\Delta D_1} \right) I,$$

where $c$ is the speed of light, $\Gamma$ is the natural linewidth, $\omega_0$ is the resonant frequency, $\Delta D_2(D_1)$ is the detuning from the $D_1(D_2)$ line.

During detection, we use a near resonant probe beam that is tuned to the $F'=4$ to $F'=5$ $D_2$ line transition so the $F'=4$, bright ground state scatters probe photons in a closed cycling transition, while the $F'=3$ state is dark to the probe and ideally scatters no light. State-dependent fluorescence is then collected and imaged onto a single photon counting module (SPCM), which is monitored in real time by our field programmable gate-array (FPGA) control system. The atom is then assigned a label bright or dark based on a discrimination threshold level of collected photon counts. The (retroreflected) probe beam is left-hand circularly polarized and a bias magnetic field (B-field) of 4 G is set counter to its propagation direction such that we drive primarily $\sigma_-$ transitions. Thus, we quickly pump the atom into the $F = 4$, $m_F = 4$ “stretched” ground state, which is protected via angular momentum conservation against depumping to the dark state through off-resonant coupling to $F' \neq 5$ states. There is also a small, transient probability of depumping during the driven random walk up to the stretched state if the atom is initialized in a different Zeeman state - see Ref. [14]. For this reason, we turn on the probe beam alongside the repump beam in the last 10 $\mu$s of optical pumping when preparing the bright state, effectively preparing the stretched state.

In an ideal scenario, we would be able to determine the state of the atom with arbitrary accuracy by collecting scattered photons until the bright and dark count distributions are sufficiently separated. However, recoil heating of the bright state during detection limits the maximum possible probe time before atom loss, and depumping from trap and probe light can lead to state information loss during detection. We study three depumping mechanisms, illustrated in Figure 1(a) off-resonant scatter of trap light, (b) probe polarization components that allow for off-resonant scatter via $F' \neq 5$, and (c) trap induced two-photon coupling of excited state hyperfine manifolds via stimulated Raman transitions. The first mechanism is the main fundamental constraint, and can be reduced by improving the photon collection rate or by decreasing the intensity of the trapping beam. The second mechanism is imposed by probe misalignment and polarization impurity, and can thus be improved with further technical capability. The third is an almost entirely geometrical problem and can be essentially eliminated by an appropriate choice of trap polarization, bias fields, probe polarization, and probe propagation direction. To our knowledge, this third information loss channel has not been previously reported for this type of detection, and turned out to be critical to achieving the state detection fidelity we report.

To guide our choice of detection parameters and predict our ultimate state-discrimination fidelity, we use a statistical model of photon collection governed by the rates of three processes: the count collection rate from a bright atom, the dark collection rate (background), and the rate of state-information loss from the prepared state of the atom (the depump rate) \[22\]. We use a detection time that is short compared to the state information loss rate, and thus we consider only the possibility of a single state change event during detection. The resulting photon collection distribution is found by marginal-
gives an effective error is the probability of assigning a bright label to an atom

For all infidelities quoted in this work, we use during the detection window.

counts from before and after a state change event that occurs when no depumping occurs, and the second term is the convolution of the photon collection rate from an atom in the prepared state, \( R_{\text{dep}} \) is the collection rate from the atom after a state changing event (i.e., collection rate from the non-prepared state), and \( R_{\text{dep}} \) is the leakage rate out of the prepared state \( \text{[23]} \). Note that the first term represents the contribution from cases where no depumping occurs, and the second term is the convolution of counts from before and after a state change event that occurs during the detection window.

Thresholding at \( m \) detected counts to assign state labels, we define the bright label probability as a function of \( F \), the prepared state: \( P_{\text{bright}}(F) = \sum_{n=m}^{\infty} P(n|F) \). The dark state readout error is the probability of assigning a bright label to an atom prepared in the dark state: \( \epsilon_{\text{dark}} = P_{\text{bright}}(F = 3) \), and similarly the bright state readout error is: \( \epsilon_{\text{bright}} = 1 - P_{\text{bright}}(F = 4) \).

For all infidelities quoted in this work, we use \( I = \frac{1}{2}(\epsilon_{\text{bright}} + \epsilon_{\text{dark}}) \), and all fidelities are \( F = 1 - I \). Numerically, we find that the optimal threshold level is \( m = 2 \) counts to discriminate between bright and dark in the range of collection and depumping rates near our experiment parameters.

The background photon collection rate is 40(1) s\(^{-1}\) and the bright atom collection rate is 2.15(1) \( \times 10^4 \) s\(^{-1}\), such that optimal error rates of \( \epsilon_{\text{bright}} = 0.005\% \) and \( \epsilon_{\text{dark}} = 0.03\% \) would be achieved in an ideal, depump-free, case at 0.59 ms of probe time. However, observed error rates are an order of magnitude higher, indicating that state information loss during detection is the dominant source of detection infidelity. We therefore focus on the state-information loss and consider the depump rate normalized by the bright collection rate \( \text{[24]} \), \( R = R_{\text{dep}}/R_{\text{bright}} \), to be a figure of merit for comparison of different mechanisms. \( R \) is the depump probability per collection event of the bright state. As a rough approximation for estimating fidelity, one can consider the probability of collecting \( m \) photons from the bright state prior to depump, \( (1 - R)^m = 1 - mR \) for \( mR \ll 1 \).

We begin our analysis of state information loss during detection by considering single photon off-resonant scattering from the trapping laser, illustrated in Fig. 1. The far detuned trapping light scatters light proportional to the intensity, and can therefore be expressed in proportion to the trap depth. Using a numeric density matrix model including all magnetic levels of the ground and first excited states \((6S_1/2 \& 6S_3/2)\), we find that the depumping rate from the trap light is \( R_{\text{dep, trap}} = 2.5 \times 10^{-7} \omega_0/h \). For a trap depth of \( \omega_0/h = 11.8(3) \) MHz, this gives an effective \( T_1 = 1/R_{\text{dep, trap}} \) of 0.34 s. To experimentally verify, we measure \( T_1 \) and find it to be 0.34(1) s. The measured value gives a normalized depump rate of \( R_{\text{trap}} = 1.5 \times 10^{-4} \) due to trap off-resonant scattering.

Trap off-resonant scattering will affect both the \( F = 3 \) and \( F = 4 \) ground states, and is a fundamental limitation of this platform. The only way to improve infidelity due to this mechanism is to collect fluorescence faster relative to the trap off-resonant scattering rate. This could be achieved by reducing the trap intensity, but decreasing the trap depth leads to a trade off in survival probability. Alternately, increasing the probe intensity could provide a higher fluorescence collection rate, but this would negatively impact the off-resonant scattering from the probe. Empirically, we also observe increased trap loss for higher probe beam intensities. Improving the photon collection efficiency would also improve the collection to depump ratio; however, doing so generally requires an optical system redesign, assuming proper operation of the current apparatus.

The probe beam may also cause state information loss due to off-resonant scatter, as illustrated in Fig. 1. Although we attempt to drive only \( \sigma_- \) transitions with the probe beam, polarization impurity or misalignment of the probe propagation direction to the B-field allows for off-resonant scatter to \( F' \neq 5 \) levels. Following the calculation in Ref. \( \text{[14]} \), we find the depump rate by summing over all dipole allowed transition scattering rates, weighted by the branching ratio \( (b_\alpha) \) of the excited state \( |(\alpha)\rangle \) to the \( F=3 \) ground state manifold.

\[
R_{\text{dep, probe}} = \sum_\alpha b_\alpha \frac{\Gamma}{2} \frac{s_\alpha}{1 + s_\alpha + (2\delta_\alpha/\Gamma)^2}
\]

The sum depends on probe polarization purity, alignment of the probe to the B-field, and intensity of the probe beam. Our measured probe polarization purity is 98.5\% before entering the vacuum chamber, and we align the B-field by scanning shim fields and maximizing the bright label probability for \( F = 4 \) after a conservatively long detection time, such that we expect an alignment tolerance of a few degrees and a resulting \( \sigma_- \) and \( \pi \)-polarization intensity fractions of \( \leq 10\% \). This corresponds to a probe off-resonant scatter rate of \( \leq 5 \) s\(^{-1}\), or \( R_{\text{probe}} \leq 2 \times 10^{-4} \); comparable to the depump probability due to the trap off resonant scatter.

Initially our measured depump rate was much higher than we would expect from the combination of contributions from trap and probe off-resonant scatter, which prompted investigation of a third mechanism. A detuned, two-photon effect, shown in Fig. 1, provides another pathway for atoms to escape the cycling transition subspace. While the ground state hyperfine splitting of \( \approx 9.2 \) GHz is sufficiently large to prevent significant A-type Raman transitions using two trap photons, the excited state hyperfine splitting is only on the order of 100 MHz. For some choices of experiment geometry, this leads to appreciable population leakage of exited state atoms to \( F' \neq 5 \) by detuned V-type Raman transitions.

As a case study for this V-type Raman mechanism, we calculate the two-photon depumping rate when the bias B-field is aligned orthogonal to the (linear) trap polarization direction,
producing $\frac{\sigma_+ + \sigma_-}{\sqrt{2}}$ trap light in the atom’s angular momentum basis. This mechanism requires and is proportional to the population of the target excited state during detection, so we naturally solve for the depump probability per detection scattering event. The single photon detuning is much larger than the Raman detuning ($|\Delta_{D_2}| \gg |\delta_{\text{fs}}|$), where the detuning values are defined in Fig. 4 so the effective two-photon Rabi rate between levels $|F' = 5, m_F = 5\rangle$ and $|F', m_F = 3\rangle$ is given by

$$\Omega_{\text{eff}} = \frac{\Omega_+ \Omega_-}{2\Delta_{D_2}}.$$  \hspace{1cm} (4)

Where $\Omega_+$ ($\Omega_-$) is the single photon Rabi rate for the $|F = 4, m_F = 4\rangle$ to $|F = 5, m_F = 5\rangle$ ($|F' = 4, m_F = 3\rangle$) leg of the Raman transition. The single photon Rabi rates of interest can be extracted from Eq. 1 using $|\Delta_{D_2}| / (2\pi) \approx 31.9$ THz and a trap depth of $U_0 / h = 11.8(3)$ MHz [26]. Since there are equal $\sigma_+$ and $\sigma_-$ components, we can easily estimate of the Rabi rate for each component by using the detunings and half the trap depth:

$$\left(\frac{\Omega_{\text{eff}}}{2\pi}\right)^2 \approx \frac{2U_0}{\hbar} \left(\frac{\Delta_{D_1}\Delta_{D_2}}{2\Delta_{D_1} + \Delta_{D_2}}\right) \approx \frac{U_0}{\hbar} \times (15.5 \text{THz}).$$  \hspace{1cm} (5)

To find the Rabi rate for the two legs of the Raman transition, $\Omega_i$ is scaled by the relative strength of each particular transition [27]. For $|F' = 5, m_F = 5\rangle$ to $|F' = 4, m_F = 3\rangle$, $\Omega_+ = \frac{3}{2}\Omega_i$ and $\Omega_- = \frac{1}{2}\Omega_i$, so that the Raman Rabi rate is $\Omega_{\text{eff}} = 21\Omega_i^2 / 160\Delta_{D_2}$. We find the population $\rho_{y'y'}$ of the $|F' = 4, m_F = 3\rangle$ excited state based on the effective Rabi rate and detuning ($\delta_{\text{fs}}^2 / 2\pi = 251\text{MHz}$ [28]) of the Raman transition to be

$$\rho_{y'y'} \approx \rho_{y'y'} \frac{\Omega_{\text{eff}}^2}{2(\delta_{\text{fs}}^2 + \Omega_{\text{eff}}^2)} \approx \rho_{y'y'} \times (5.6 \times 10^{-6}).$$  \hspace{1cm} (6)

The branching ratio of the $F' = 4$ manifold to $F = 3$ ground state is $5/12$ (see Table I), so the overall probability of depumping per resonant scattering event for this transition is

$$p_{\text{depump,}y'y'} = b_{y'y'} \frac{\rho_{y'y'}}{\rho_{y'y'}} \approx 2.3 \times 10^{-6}. \hspace{1cm} (7)$$

A similar calculation can be carried out for $|F' = 3, m_F = 3\rangle$ on the $\sigma_-$ leg, for which $p_{\text{depump,}y'y'} \approx 3.6 \times 10^{-6}$. Combining these two depump probabilities, we have a $6 \times 10^{-6}$ chance of depump for each resonant scattering event with this configuration, or $R_{\pi} = 1.6 \times 10^{-3}$ after accounting for our collection efficiency (CE) of 0.37%, which is a measured value that includes transmission of our optical system and the quantum efficiency of the detector. Since this value is an order of magnitude larger than those of either single photon off-resonant depumping rates, mitigation of this mechanism yields significant gains for the bright state readout fidelity.

To verify this newly identified mechanism, we study the bright state count histogram for two experiment geometries. We orient the trap polarization either parallel or orthogonal to the B-field/probe-defined quantization axis generating either $\pi$ or $\sigma_-$ polarization components of the trapping beam [29]. With the $\sigma_-$ polarized trap light, we expect to see depumping from the two-photon pathway as calculated, however for the $\pi$-polarized trap configuration, there is no dipole-allowed coupling from the excited state to the ground state caused by the trap light, and thus we expect this mechanism to be absent (barring trap polarization misalignment and polarization distortion effects from focusing).

As shown in Fig. 5 when we run the experiment in these two configurations, we find that the histogram of collected light from a bright atom in a $\pi$-polarized trap is significantly
FIG. 3: We measure the wait time distribution for the bright and dark state using adaptive detection. 5 µs probe pulses are applied until the threshold level of counts are recorded or the maximum possible detection time of 500 µs is reached. The time distribution of the bright state (red bars) closely matches an Erlang distribution (black line), which is the expected distribution for the depump-free case. The discrepancy between the chance that bright atoms that reach the maximum detection time (0.15±0.06%) and the Erlang distribution at max time reveals the probability that the atoms become dark before the threshold level of photons is collected. The time distribution for atoms prepared in the dark state (blue bars) is nearly singular at the maximum time bin. The combined effect of background counts on the detector and dark atoms becoming bright allows for small, but non-zero, probability of reaching threshold counts and stopping detection early when preparing in the dark state. This data was produced by recording the number of pulses applied each time when collecting data in Figure 4.

We measure the wait time distribution for the bright atom in a σ-polarized trap, implying a greater overall depump rate in the σ-polarized trap. A fit to Eq. (2) yields depumping probabilities per scattering event of 50(5) × 10^{-6} and 7(4) × 10^{-6} for the σ and π configurations respectively. Since the other probe and trap parameters are held fixed, we attribute the difference in the depump rates between the two configurations to the off-resonant, V-type Raman transition. This data was taken prior to a significant optimization of the probe beam optics (alignment and beam quality), and we attribute the remaining depump rate in the π configuration mostly to probe impurities.

Achieving excellent state readout fidelity is crucial for this platform; however, from a practical standpoint atom retention is equally important, and can be a challenging problem to solve. Photon recoil from the near resonant probe beam causes a bright atom to undergo a random walk in momentum space, eventually allowing the atom to gain sufficient energy to overcome the trap depth \((U_0/k_B = 0.57(1) \text{mK})\) and escape. Other heating mechanisms, such as dipole-force fluctuation heating \([30]\) and ODT intensity fluctuation heating \([31]\) are neglected here due to their small contributions on this experiment, as observed by small AC Stark shift on the probe transition \([32]\) and long (seconds scale), trap power independent atom storage time respectively.

In a standard, single-pulse detection scheme, the mean number of counts collected from the bright state must be significantly higher than the threshold level in order to collect above-threshold counts from the bright state with high probability. This poses a problem when the recoil heating becomes comparable to the trap depth. For example, to achieve \(\geq 99.9\%\) bright label probability from a Poisson count distribution by thresholding at 2 or more counts, the mean would need to be \(\geq 9.23\) counts. Dividing by our collection efficiency of \(CE = 0.37\%\) gives a requirement of \(N_{\text{cut}} \approx 2.500\) photons scattered on average. If we approximate the expected heating from photon recoil with the free atom case, a resonant scattering event adds \(\Delta E/k_B = 2T_{\text{rec}} \approx 2 \times 0.2 \mu K\) per recoil event \([28, 30, 33]\), such that the mean energy gain after single pulse detection would be 1 mK, nearly twice the trap depth.

However, if state labelling is done entirely with thresholding, any extra light collected beyond the threshold level provides no additional information. Therefore, we can avoid much of the unnecessary recoil heating if we turn off the detection light as soon as we have reached the threshold level. In the ideal case of stopping the instant the threshold photon is received, the mean number of photons collected for a depump-free bright atom would simply be the threshold level, and thus the average temperature gain would be roughly a factor of \(2/9.23\) times lower than the constant probe time case, keeping the average energy gain below the trap depth.

We refer to this protocol as “adaptive detection,” and note that similar protocols were used on neutral atoms in Ref. \([34]\) and on trapped ions in Refs. \([35–39]\). We implement such a scheme by applying a series of short (5 µs) pulses until either the threshold number (2 counts) is detected, or we reach a total integrated probe duration of \(t_d = 500\) µs. This maps the count distribution in Eq. (2) from the bright state onto a nearly Erlang wait time \((t)\) distribution with shape parameter equal to the detection threshold \((m = 2): f(t) \approx R^2_{\text{bright}}e^{−R_{\text{bright}}t}\) shown in Fig. 3.

We measure our detection fidelity by alternately preparing the atom in each state 100,000 times and recording the number of times that the measured state matches the prepared state, see Fig. 4. We collect the data in 1,000 batches of 100 shots, alternating between bright and dark state preparation for each batch. A total of 152 atoms prepared in the bright state are read out as dark and 159 atoms prepared in the dark state are labelled bright, giving a bright error of \(\epsilon_{\text{bright}} = 0.15±0.03\%\) and a dark error of \(\epsilon_{\text{dark}} = 0.16±0.03\%\), where the uncertainty is a Wilson score interval. This an infidelity of 0.16% or a fidelity of 99.84\%. This number does not correct for state
FIG. 4: The normalized histogram of collected photons from the bright and dark state demonstrates 0.16 ± 0.03% detection infidelity. Single atoms are prepared in either the dark (blue) or bright (red) state and read out by illuminating the atom with the D2 line probe and collecting scattered photons. When two or more photons are collected, the atom is labelled bright; when fewer than two photons are collected, the atom is labelled dark. Atoms prepared in the bright state are erroneously labelled dark in 0.15 ± 0.02% of experiments, and atoms prepared in the dark state are erroneously labelled bright in 0.16 ± 0.03% of experiments. Data presented in this histogram was collected in 1,000 batches of 100 shots, alternating between dark and bright state preparation, for a total of 100,000 shots for each state.

preparation errors, although we expect our state preparation error to be small compared to the readout error. We then extract the detection-driven loss probability by subtracting the ratio of the probability of passing the presence check when preparing in the bright and dark state from unity, and find 2.6(2)% loss of the bright state due to detection. We then decrease our trap depth from 11.8(3) MHz to 5.9(2) MHz to reduce the off resonant scattering from the trap light, at the cost of reduced survival probability of the bright state. This allows us to reach a detection fidelity of 99.89 ± 0.03% with 14.1(3)% loss of the bright state due to detection.

Casting the bright state readout error as the total depump probability per collection event, we find $R = 1 - \sqrt{1 - \epsilon_{\text{bright}}} = 7.5 \times 10^{-4}$. This is the same order of magnitude as the combined estimate for trap and probe off-resonant scatter. The source of the remaining discrepancy in the depump rate relative to the estimates for off-resonant scatter is a subject of ongoing study, potentially due to an overestimation of the alignment precision of the probe, imperfect state preparation, or residual $\sigma$-polarized components of the trap light due to misalignment of the trap polarization direction to the B-field or polarization distortion effects due to focusing of the trap beam.

We note that further improvement of the detection fidelity and atom retention should be achievable by improving the photon collection efficiency. Our measured collection efficiency of 0.37% is relatively poor compared to other works [13,15] and we expect that the detection fidelity should improve approximately linearly with the ratio of collected photons to depump rate. We also note that these results ease vacuum system requirements for high fidelity detection of alkali atoms in optical tweezers, since detection fidelity is independent of the background atom loss rate. To achieve comparable fidelities with a push out method, the average atom loss rate due to background gas collisions (or indeed any source other than the pushout beam) must be comparable to the dark readout error reported here, since a dark atom lost due to background is indistinguishable from a bright atom lost to the pushout beam in such schemes. Finally, we note that our result obviates the need for toggling the trap and probe beams, as we keep the trap beam on at all times and still achieve high fidelity. Use of a magic wavelength trap is important for our ability to keep the trap on, as we avoid dipole force fluctuation (DFF) heating [30]. Toggling the trap and probe beam could be an alternate strategy to mitigate both the two-photon depumping mechanism and DFF heating, at the cost of potentially introducing an extra heating mechanism from repeated kicks.

Ultimately, the fidelity reported here is the result of study of the relevant atom-photon interaction physics, and mitigation of a previously unreported depumping mechanism under simultaneous trap and probe illumination. We find that the remaining infidelity from off-resonant scattering of trap and probe light is sufficiently low to enable an order of magnitude improvement over previously published low-loss readout of alkali atoms in optical tweezers. In conjunction with the high atom retention enabled by adaptive detection, this result alludes to promise of near-term detection suitable for fault tolerant operation and the possibility of non-disruptive mid circuit measurement for error detecting/correcting algorithms. This work represents an important step towards building scalable, high-performance quantum information processors and quantum sensors out of alkali-metal atoms trapped in optical tweezer traps.

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$|\alpha|^2 P(F, m_F')$ Polarization $\delta \alpha / 2 \pi$ (MHz) $b_{sat}$ $I_{sat, \alpha} / I_{sat}$

|   |   |   |   |
|---|---|---|---|
| 4, 4 | $\pi$ | 251 | 5/12 | 2.14 |
| 4, 3 | $\sigma_{-}$ | 251 | 5/12 | 8.57 |
| 3, 3 | $\sigma_{-}$ | 452 | 3/4 | 5.14 |

TABLE I: Summary of relevant parameters for calculation of probe-driven off-resonant coupling to the $|F = 4, m_F = 4\rangle$ ground state. $|\alpha\rangle$ is the final state, $\delta \alpha / 2 \pi$ is the detuning of the probe beam from that transition, $b_{sat}$ is the branching ratio to the $F=3$ ground state manifold, and $I_{sat, \alpha} / I_{sat}$ is the saturation intensity of the transition, normalized to that of $|F = 4, m_F = 4\rangle$ to $|F = 5, m_F' = 5\rangle$ transition, which is the resonant transition for probe beam. The probe is mostly $\sigma_{-}$-polarized, pumping population into $|F = 4, m_F = 4\rangle$, so only excited states with non-zero $b_{sat}$ are considered. Detunings are taken from Ref [28], and branching ratios and saturation intensities are calculated from transition strengths given in Ref [27].