Magneto-optical trap of a Group III atom

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We realize the first magneto-optical trap of an atom in main group III of the Periodic Table. Our atom of choice (indium) does not have a transition out of its ground state suitable for laser cooling; therefore, laser cooling is performed on the \( |5P_{3/2}, F = 6 \rangle \rightarrow |5D_{5/2}, F = 7 \rangle \) transition, where \( |5P_{3/2}, F = 6 \rangle \) is a long-lived metastable state. Optimization of our trap parameters results in atoms numbers as large as \( 5 \times 10^6 \) atoms with temperatures of order 1 mK. Additionally, through trap decay measurements, we infer a one-body trap lifetime of 12.3 s. This lifetime is consistent with background gas collisions and indicates that our repumpers have closed all leakage pathways. We also infer a two-body loss rate of \( 1.6 \times 10^{-11} \) cm\(^3\)/s, which is comparable to those measured in alkali atoms. The techniques demonstrated in this work can be straightforwardly applied to other group III atoms, and our results pave the way for realizing quantum degenerate gases of these particles.

The magneto-optical trap (MOT) is the workhorse of ultracold physics. This technique has enabled the realization of quantum degenerate gases [1], quantum simulators [2–4], exotic quantum matter [5], neutral atom quantum processors [6–9], record-accurate atomic clocks [10], and they also allow for precise probes of atoms’ internal states and their interactions accurately atomic clocks [10]. Therefore, triels could be probed with the precision of an atomic clock and offer the many-body control of alkalis. Furthermore, unlike alkali earths, the triel atoms thallium and indium have narrow linewidth electronic transitions that are amenable to stable laser technology. These transitions (which are not found in alkalis) are used in highly accurate atomic clocks [10], and they also allow for precise probes of atoms’ internal states and their interactions [13–15]. Therefore, triels could be probed with the precision of an atomic clock while offering the many-body control of alkalis. Furthermore, unlike the \( S \)-orbital ground states of alkalis and alkaline earths or the high angular momentum ground states of popular lanthanides like erbium and dysprosium, the \( P \)-orbital ground states of triels distinguish themselves as intermediate cases. Low-temperature scattering between such states has not been studied, and these interactions could prove as interesting and surprising as scattering measurements with lanthanides [16–19].

The main challenge of realizing a triel MOT is that these atoms do not contain cycling transitions amenable to laser cooling out of their \( F_{1/2} \) ground states. Earlier attempts to transversely cool beams of the triel atom indium concluded that cooling on the \( 5P_{1/2} \rightarrow 6S_{1/2} \) transition is inefficient even with repumping [20]. However, the long lived metastable \( P_{3/2} \) states in triel atoms offer a \( P_{3/2} \rightarrow D_{5/2} \) cycling transition that is suitable for laser cooling. Using this transition, transverse cooling was observed in the triel atoms Al [21], Ga [22], and In [23], and our team recently realized an indium Zeeman slower [24]. Laser cooling of the triel Tl has also been proposed with this transition [25]. Although most MOTs are formed using transitions out of atomic ground states, a MOT based on a transition out of an optically pumped metastable electronic state (in the lanthanide europium) [26] has been demonstrated. However, a MOT has never been realized with a triel atom.

In this Letter, we demonstrate the first magneto-optical trap (MOT) of a group-III atom using \(^{115}\)In. This species is a boson with a nuclear spin of \( I = 9/2 \) and an isotopic abundance of 95.7%. The \(^{115}\)In energy level diagram and relevant transitions for this work are shown in Fig. 1(a). For efficient laser cooling, we use the 326 nm \( |5P_{3/2}, F = 6 \rangle \rightarrow |5D_{5/2}, F = 7 \rangle \) transition. Possible mechanisms that cause population to leak out of this transition’s states are the \( 5P_{3/2} \rightarrow 5P_{1/2} \) lifetime (predicted to be \( \sim 10 \) s [31]), the \( 5D_{5/2} \rightarrow 6P_{3/2} \rightarrow 6S_{1/2} \) decay pathway (which occurs with a branching ratio of order \( 10^{-4} \) [30]), and off-resonant driving of the \( |5P_{3/2}, F = 6 \rangle \rightarrow |5D_{5/2}, F = 6 \rangle \) transition. To close these leaks, we repump population into the \( |5F_{3/2}, F = 6 \rangle \) cooling state with two 410 nm repumpers (which address the \( |5P_{1/2}, F = 4, 5 \rangle \rightarrow |6S_{1/2}, F = 5 \rangle \) transitions) and two 451 nm repumpers (which address
An indium atomic beam is produced by an effusion cell operating at 800 °C, resulting in a $10^{-3}$ Torr vapor pressure inside the crucible. The output of the effusion cell is collimated with 200 microchannels that are 1 cm in length and 200 μm in diameter. Atoms that emerge from the cell are pumped out of the $5P_{3/2}$ ground state with a pair of 410 nm lasers, which results in ample population in the $|5P_{3/2}, F = 6\rangle$ cooling state due to a favorable branching ratio. Additionally, we use two 451 nm lasers to pump the remaining population that decays into $|5P_{3/2}, F = 4, 5\rangle$ into $|5P_{3/2}, F = 6\rangle$. The atomic beam is then slowed to a 70 m/s velocity with a transverse field permanent magnet Zeeman slower [24].

The slowed atomic beam enters a vacuum chamber at a pressure at the low $10^{-10}$ Torr level. A pair of coils in an anti-Helmholtz configuration generate a quadrupole magnetic field that has an axial field gradient $\partial B/\partial z$ up to 55 G/cm at the center of the chamber [Fig. 1(b)]. Three pairs of orthogonal counter-propagating 325.7 nm cooling beams, which have a $1/e^2$ diameter of 12 mm, are circularly polarized in the traditional $\sigma^+ - \sigma^-$ configuration. In addition, the four repumpers, which also have a $1/e^2$ diameter of 12 mm, are coaligned with the $z$ direction MOT beam. The cooling laser is generated by frequency quadrupling the output of a 1302.8 nm external-cavity-diode-laser-seeded Raman fiber amplifier, resulting in hundreds of milliwatts of useful 325.7 nm laser power. The 325.7 nm laser is frequency stabilized to a calibrated, low-drift HighFinesse wavemeter. Meanwhile the 4 repumper beams are generated by external cavity diode lasers stabilized to indium hollow cathode lamps with Doppler-free spectroscopy.

With these lasers in place, we observe a bright MOT signal. MOT fluorescence is collected with a 2$f$ imaging system focused onto an electron multiplying CCD (EM-CCD) camera. The number of atoms $N$ in the MOT is inferred from a fluorescence image as [32]

$$N = \frac{8\pi N_c}{\eta \Omega t_{ex} s_0} \left( 1 + s_0 + \frac{4\Delta^2}{\Gamma^2} \right),$$

where $N_c$ is the total number of fluorescence counts in a MOT image, $\Omega$ is the solid angle subtended by the imaging system, $\eta$ is the imaging system’s quantum efficiency, $t_{ex}$ is the exposure time of the camera, $\Delta$ is the MOT detuning (in rad/s), and $\Gamma$ is the natural linewidth (in rad/s) of the cooling transition. The saturation parameter of the MOT lasers is $s_0 = I/I_{sat}$, where $I$ is the total intensity of all 6 MOT laser beams.
and $I_{sat} = 78.3 \text{ mW/cm}^2$ is the saturation intensity of the cooling transition. The multiplicity factor $\mathcal{M} = 3(2F + 1)/(2F + 3) = 2.6$ (where $F$ is associated with the lower-energy cooling state) accounts for the fact that the cooling transition is not a two-level system; rather, the MOT acts on the Zeeman sublevels of each cooling state, and population is assumed to be distributed roughly evenly among these sublevels [26, 32–38].

Although Eq. 1 ignores the magnetic field, the field’s fractional contribution to the atom number calculation is [32]

$$\frac{\Delta N}{N} = \left(\frac{2\mu}{\Gamma} \frac{\partial B}{\partial z}\right)^2 \frac{\sigma^2 + \frac{1}{2} \sigma^2_{y,z}}{1 + s_0 + 4 \Delta I / \Gamma}. \quad (2)$$

Here $\Delta N$ is the change in the computed atom number due to the magnetic field, $\mu$ is the effective magnetic moment, $\partial B/\partial z$ is the $z$ derivative of the magnetic field evaluated at the trap center (taken to be the origin), and $\sigma_{x,y,z}$ are the RMS radii of the Gaussian atomic density distribution. We estimate that including the magnetic field results in a 3.5% correction to Eq. 1, so we neglect this.

![Figure 2: Optimization of MOT atom number. We vary the MOT detuning $\Delta$, the MOT laser intensity (expressed in the Figure as the saturation parameter $s_0 = I/\mathcal{M}I_{sat}$), and the magnetic field gradient $\partial B/\partial z$ at the trap center projected along the $z$ axis. The measurement is taken by fixing the detuning and then varying $s_0$ and $\partial B/\partial z$ until the atom number maximizes. Each bar represents the best atom number observed for a given detuning. The trapped atom number fluctuations are estimated from the standard deviation of multiple atom number measurements for the same MOT parameters.](image)

To optimize the total atom number, we systematically vary $\Delta$. Once a value of $\Delta$ is fixed, we vary the per-beam MOT laser power $P$ and $\partial B/\partial z$ until the atom number is maximized. We find a global maximum of the atom number when $\Delta/\Gamma = -1.6$, $\partial B/\partial z = 36 \text{ G/cm}$, and $P = 24.5 \text{ mW}$, the latter of which corresponds to $s_0 = 1.3$ (Fig. 2). With these values, we observe a MOT with $5 \times 10^8$ atoms, which competitive with the large atom numbers achieved with alkali [39] and alkaline earth [40] MOTs. The uncertainty due to random fluctuations in $N$ is obtained by measuring the atom number several times for a fixed set of MOT parameters. The standard deviation in these measurements divided by their mean is treated as the fractional uncertainty in $N$ and applied to the remainder of the optimization data (Fig. 2).

With the MOT parameters fixed at their optimum values, we measure the temperature of the trapped gas with a time-of-flight measurement. We observe the vertical (horizontal) temperature to be $3.6 \text{ mK}$ ($4.5 \text{ mK}$) for optimal MOT parameters, while the predicted temperature from Doppler theory is $1 \text{ mK}$. MOT temperatures above the Doppler value have been observed in multiple experiments [41–47].

An important characterization of a MOT is its one- and two-body loss rates. A small one-body lifetime can indicate that our repumping scheme has not closed all population leaks out of the cooling states. This is especially concern since our cooling is based on a metastable state, which results in many more levels lower in energy than the cooling transition (and therefore more decay pathways) compared to alkali and alkaline earth MOTs that cool out of a ground state hyperfine manifold. In fact, another realization of a MOT using an optically pumped metastable electronic transition found a small one-body lifetime and identified uncontrolled population leaks as the cause [26]. Additionally, a measurement of the two-body rate can determine the severity of unwanted effects like light-assisted collisions.

We measure these rates by disrupting MOT loading and observing the resulting atom number decay (Fig.3). The MOT loading rate is changed by shutting off the Zeeman slower laser, which results in a greatly reduced steady-state MOT number. Our decay data is taken with the optimal MOT parameters (found above) as well as $7 \text{ mW}$ of combined 410 nm laser repumper power and $7 \text{ mW}$ of combined 451 nm laser repumper power.

The decay of the trapped atom number $N(t)$ in the MOT can be modeled by a rate equation,

$$\frac{dN}{dt} = L - \frac{N}{\tau} - \beta \int n^2(\vec{r}, t)d^3\vec{r}, \quad (3)$$

where $L$ is the MOT loading rate (in this case, the rate without the Zeeman slower laser), $\tau$ is the MOT one-body lifetime, and $\beta$ is the two-body rate constant. When the MOT has a Gaussian atomic density distribution, we can simplify Eq. 3 by rewriting the two-body loss term as $-\Gamma_2 N^2$, where $\Gamma_2 = \beta/(8\pi^3/2\sigma_x\sigma_y\sigma_z)$. The decay of the trapped atom number $N(t)$ is therefore given by

$$N(t) = \frac{2L\tau - (1 - \kappa)N_0}{1 + \kappa + 2N_0\Gamma_2\tau} e^{-\kappa t/\tau}, \quad (4)$$
where \( \kappa = \sqrt{1 + \frac{4L}{\Gamma^2 \tau^2}} \) and \( N_0 \) is the atom number at \( t = 0 \). The MOT lifetime and two-body loss rate can be extracted from fitting the decay data to Eq.4. For the one-body lifetime, we obtain a value of \( \tau = 12.3(3) \) s. This number is consistent with decay due to collisions with the background gas, which is the technical limit of the MOT lifetime. Meanwhile, the fitted two-body rate constant is \( \beta = 1.6(5) \times 10^{-11} \) cm\(^3\)/s. We have observed that the fitted value of \( \beta \) varies by 30\% between different runs of our experiment, and this variation is reflected in our quoted uncertainty for \( \beta \). For comparison, two-body rate constants measured in well-functioning MOTs of alkali atoms are found to be \( 4 \times 10^{-11} \) cm\(^3\)/s for \( ^{23}\text{Na} \), \( 5.8 \times 10^{-12} \) cm\(^3\)/s for \( ^{87}\text{Rb} \), and \( 7.6 \times 10^{-11} \) cm\(^3\)/s for \( ^{133}\text{Cs} \) [50].

To further study repumping efficiency, we vary repumper powers and measure the effect on the MOT one-body lifetime and steady-state atom number. Data is collected by varying the combined power of both 410 nm (451 nm) repumpers while the power of the 451 nm (410 nm) repumpers is fixed at its maximum value. Below 4 mW (1 mW) of 410 nm (451 nm) laser light, we observe that the one-body lifetime and steady-state atom number depends strongly on repumper power (Fig. 4). Above this regime, the repumping transitions are saturated, and the MOT has a large atom number and a long one-body lifetime. Comparing the 410 nm and 451 nm repumping data, it is clear that 451 nm repumping saturates more easily. This is because the 410 nm lasers are somewhat effective at populating \([5P_{3/2}, F = 6]\) on their own due to a 60\% branching ratio for decay into this state.

We compare this result to MOT realizations that are reportedly limited by repumper power or population leakage pathways that are not fully closed [26, 36, 51–54]. In these cases, one-body lifetimes are an order-of-magnitude shorter than in our apparatus. The data reported here suggest that population leaks are well closed by our repumpers.

In conclusion, we demonstrate a magneto-optical trap of a group-III atom. Over \( 5 \times 10^8 \) atoms are loaded to the MOT from a Zeeman slowed atomic beam. With ample repumper power, measurements of MOT decay imply a one-body lifetime of \( \tau = 12.3(3) \) s. This lifetime as well as repumper characterization confirms that all leakage pathways are sufficiently closed. Meanwhile, we observed a two-body rate of \( \beta = 1.6(5) \times 10^{-11} \) cm\(^3\)/s. Further improvements are possible by increasing the atomic density using dark spot MOT or compressed MOT techniques, and lower temperatures can be achieved with Siyaphus cooling. Techniques demonstrated in this work can be extended to all atoms with similar atomic structure, such as the group-III atoms Tl, Al, Ga, and B. Furthermore, our work paves the way of exploring group-III atoms in the quantum degenerate regime.
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I. INFERRING THE ATOM NUMBER FROM A MOT FLUORESCENCE MEASUREMENT

We consider a small volume of space \( dV \) located at coordinates \( x, y, \) and \( z \). The volume is small enough that the external magnetic field gradient does not vary significantly within it and can therefore be taken as constant. The rate of photons emitted from this volume is

\[
d\Gamma_{\text{ph}} = \Gamma_{\text{sc}} dN,
\]

where \( dN \) is the atom number within \( dV \), and the single-atom scattering rate \( \Gamma_{\text{sc}} \) is

\[
\Gamma_{\text{sc}} = \frac{\Gamma}{2} \frac{1 + s_0^2 + 4\Gamma^2(\Delta + \mu_B)^2}{1 + s_0 + \frac{4\Gamma^2}{\chi^2}(\Delta + \mu_B)^2}.
\]

Here \( \Gamma \) is the spontaneous emission rate from the MOT transition, \( s_0 \) is the saturation parameter, \( \Delta \) is the MOT detuning, and \( B \) is the MOT’s external magnetic field. The effective magnetic moment of the MOT transition is

\[
\mu = (g_e m_e - g_g m_g) \mu_B,
\]

where \( g_e \) (\( g_g \)) is the excited (ground) state \( g \)-factor, \( m_e \) (\( m_g \)) is the excited (ground) state magnetic quantum number, and \( \mu_B \) is the Bohr magneton.

The quadrupole field \( B \) can be approximated near the origin as

\[
B \simeq \beta_x x + \beta_y y + \beta_z z.
\]

Here \( \beta_i \) is defined as

\[
\beta_i = \frac{\partial B}{\partial i} \bigg|_{x=y=z=0},
\]

with \( i = x, y, z \).

The rate of photons emitted from \( dV \) that make it to the camera is

\[
d\Gamma_{\text{cam}} = \frac{\Omega}{4\pi} d\Gamma_{\text{ph}} = \frac{\Omega}{4\pi} \Gamma_{\text{sc}} dN,
\]

where \( \Omega \) is the solid angle subtended by the imaging system. One useful expression for this analysis is the density of camera counts \( n_c \), which can be obtained by fitting a picture of the MOT fluorescence with a Gaussian, resulting in

\[
n_c = \frac{N_c}{(2\pi)^{3/2}\sigma_x\sigma_y\sigma_z} e^{-x^2/2\sigma_x^2} e^{-y^2/2\sigma_y^2} e^{-z^2/2\sigma_z^2}.
\]

Here \( N_c \) is the total integrated camera counts, and \( \sigma_x, \sigma_y, \) and \( \sigma_z \) are the fitted RMS widths of the MOT (usually only two of these widths can be fit and the other is inferred by symmetry). Using this expression, the number of camera counts due to the fluorescence originating from \( dV \) is

\[
n_c dV = \frac{\Omega}{4\pi} \Gamma_{\text{sc}} \eta t_{ex} dN,
\]

where \( \eta \) is the camera’s quantum efficiency and \( t_{ex} \) is its exposure time. Or in terms of densities,
\[ n_c = \frac{\Omega}{4\pi} \Gamma_{sc} \eta_{ex} n, \]  

where \( n \) is the atom number density.

The total atom number is therefore

\[ N = \int n \, dV \]  
\[ = \frac{4\pi}{\Omega \eta_{ex}} \int \Gamma_{sc} n_c \, dV \]  
\[ = \frac{8\pi}{\Omega \Gamma \eta_{ex} s_0} \left[ \left( 1 + s_0 + \frac{4\Delta^2}{\Gamma^2} \right) N_c + \frac{4}{\Gamma^2} \int (2\Delta \mu_B + \mu^2 B^2) n_c \, dV \right]. \]

Note since \( B \) is an odd function in its spatial coordinates whereas \( n_c \) is an even one, the first term in the integral of Eqn. 11 vanishes. The main text (and several other references) estimates the atom number by ignoring the magnetic field, resulting in the approximate number expression

\[ N_{\text{approx}} = \frac{8\pi N_c}{\Omega \Gamma \eta_{ex} s_0} \left( 1 + s_0 + \frac{4\Delta^2}{\Gamma^2} \right). \]

The correction to \( N_{\text{approx}} \) due to the magnetic field is therefore

\[ N - N_{\text{approx}} = \frac{32\pi \mu^2}{\Omega \Gamma^3 \eta_{ex} s_0} \int B^2 n_c \, dV \]  
\[ = \frac{32\pi \mu^2 N_c}{\Omega \Gamma^3 \eta_{ex} s_0} \left( \beta_x^2 \sigma_x^2 + \beta_y^2 \sigma_y^2 + \beta_z^2 \sigma_z^2 \right). \]

Using the fact that \( \sigma_x = \sigma_y \) and \( \beta_x = \beta_y \) (due to the symmetry of the anti-Helmholtz field) as well as \( \beta_x + \beta_y + \beta_z = 0 \) (Gauss’ law for magnetism), we can write the fractional correction to the approximate atom number calculation as

\[ \frac{N - N_{\text{approx}}}{N_{\text{approx}}} = \frac{4\mu^2 \beta_z^2}{\Gamma^2} \frac{\sigma_z^2 + \frac{1}{3} \sigma_x^2}{1 + s_0 + \frac{4\Delta^2}{\Gamma^2}}. \]

For the optimal numbers found in the main text, \( \mu = \mu_B, \beta_z = 36 \text{ G/cm}, s_0 = 1.3, \text{ and } \Delta/\Gamma = -1.6. \) Also, \( \sigma_x = 1.3 \text{ mm} \) and \( \sigma_z = 1 \text{ mm}. \) This results in a correction of 3.5% to the atom numbers quoted in the main text. All numbers in the main text neglect this correction.

II. SATURATION INTENSITY MULTIPLICITY FACTOR

Consider an atomic transition driven by a laser field. The saturation intensity is

\[ I_{\text{sat}} = \frac{\hbar \omega}{2 \sigma \tau}, \]

where \( \omega \) is the laser frequency, \( \sigma \) is the photon scattering cross section, and \( \tau \) is the transition lifetime. For a two-level system, the photon scattering cross section is

\[ \sigma = \frac{3\lambda^2}{2\pi} \frac{1}{1 + \frac{4\Delta^2}{\Gamma^2}}. \]

Here, \( \lambda \) is the photon wavelength. On resonance, this results in the familiar expression for the saturation intensity
\[ I_{\text{sat}} = \frac{\pi hc}{2\lambda^3 \tau} \]  

(18)

When the laser is driving a transition between two atomic hyperfine states, each with a manifold of Zeeman sublevels, Eq. 17 is no longer valid. However, it can be replaced with the effective cross section [1]

\[ \sigma_{\text{eff}} = \frac{3\lambda^2}{2\pi} \frac{\mathcal{M}}{1 + \frac{4\lambda^2}{\Delta^2}} \]

(19)

\[ = \mathcal{M}\sigma, \]

(20)

where the multiplicity factor \( \mathcal{M} \) is

\[ \mathcal{M} = \frac{3(2F + 1)}{2F + 3}. \]

(21)

Therefore, in the presence of Zeeman sublevels, the resonant saturation intensity is \( \mathcal{M}I_{\text{sat}} \), where \( I_{\text{sat}} \) is the two-level saturation intensity of Eq. 18.

[1] C. A. Muller, C. Miniatura, D. Wilkowski, R. Kaiser, and D. Delande, Multiple scattering of photons by atomic hyperfine multiplets, Phys. Rev. A 72, 053405 (2005).