Recoil-sensitive lithium interferometer without a subrecoil sample

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We report simultaneous conjugate Ramsey-Bordé interferometers with a sample of low-mass (lithium-7) atoms at 50 times the recoil temperature. We optically pump the atoms to a magnetically insensitive state using the $2S_{1/2} - 2P_{3/2}$ line. Fast stimulated Raman beam splitters address a broad velocity class and unavoidably drive two conjugate interferometers that overlap spatially. We show that detecting the summed interference signals of both interferometers, using state labeling, allows recoil measurements and suppression of phase noise from vibrations. The use of “warm” atoms allows for simple, efficient, and high-flux atom sources and broadens the applicability of recoil-sensitive interferometry to particles that remain difficult to trap and cool.

In a light-pulse atom interferometer, laser pulses with wavenumber $k$ direct matter waves along a superposition of trajectories and recombine them to reveal the phase difference between paths $|1|$. They are used for inertial sensing $[2, 3]$, gravity gradiometry $[4]$ and tests of fundamental physics $[3-15]$. Ramsey-Bordé interferometers, in particular, measure the mass $m$ of an atom through the kinetic energy $\hbar \omega_r = \hbar^2 k^2 / (2m)$ it gains after recoiling from the interaction with a photon ($\hbar$ is the reduced Planck constant). They can help redefine the kilogram $[16, 17]$ and determine the fine-structure constant $[18–22]$, thereby testing the Standard Model $[23, 24]$. The recoil frequency $\omega_r$, and therefore the signal, scales inversely with mass. Light atomic species have been used in supersonic atomic-beam interferometers $[25, 26]$, but remain difficult to cool below the recoil temperature $T_r$ where the average thermal speed equals the recoil velocity. This makes it impossible to spatially resolve the interferometer outputs, which is required for direct rejection of common-mode inertial signals with phase extraction methods $[27, 29]$.

Here, we demonstrate recoil-sensitive interferometry with a sample of lithium-7 atoms well above the atomic recoil temperature (50$T_r$), the first interferometer with laser-cooled lithium atoms or any atom lighter than sodium-23 $[30]$. Fast Raman transitions $[31]$ ($\tau_{\pi/2} = 160$ ns) address the ensemble’s large Doppler spread and simultaneously drive overlapped conjugate Ramsey-Bordé interferometers. Superimposing simultaneous conjugate interferometers suppresses effects from two-photon detuning and unwanted inertial signals, such as vibrations. Our measurement sensitivity benefits from lithium’s high recoil frequency of $\omega_r = 2\pi \times 63$ kHz (compared to $2\pi \times 2$ kHz for cesium) and the absence of time-consuming additional cooling $[32]$ or lossy velocity selection $[33]$ steps that reduce sample size and precision. The lithium isotopes present an attractive pair for testing Einstein’s Equivalence Principle using light-pulse atom interferometry $[34]$. This work broadens the applicability of recoil-sensitive interferometry to other particles: electrons $[35]$, for example, boast GHz-recoil frequencies and would enable observation of relativistic effects $[16, 36]$.

Figure 1(a) shows the trajectories of an atom in a
Ramsey-Bordé sequence. Atom-light interactions are used to split, redirect and interfere the atomic matter waves. The Ramsey-Bordé sequence consists of four $\pi/2$ (beam-splitter) pulses, so that the lowest interferometer arm remains stationary. The outputs of the second pulse that do not contribute to $A^-$ and $B^-$ may form another conjugate (upper) interferometer with final outputs $A^+$ and $B^+$. In each interferometer, the probability of detecting the atoms at one output depends on the phase difference between the arms of the interferometer, which we denote $\Delta \phi^-$ (lower interferometer) and $\Delta \phi^+$ (upper interferometer). Using standard methods\cite{37}, $\Delta \phi^+$ is calculated to second order in $T$ as

$$\Delta \phi^\pm = \pm 8\omega_r T - 2ka_z(T + T') - 2\delta T \quad (1)$$

The first term arises from the atomic kinetic energy, the second from any acceleration $a_z$ (such as gravity and vibrations) along the laser beam axis, where the average wave number of the counter-propagating beams is $k = (k_1 + k_2)/2$, and the third from the detuning of the laser frequencies from two-photon resonance in the absence of AC stark shifts, $\delta = \omega_1 - \omega_2 - (\omega_A - \omega_B)\ [41].$

The interferometers in Fig. I(a) share the first and second beam-splitter pulses. For the third and fourth pulse, the lower interferometer requires a transition coupling $|F = 2, p = 0\rangle \rightarrow |F = 1, p = -2\hbar k\rangle$, and the upper interferometer requires coupling $|F = 1, p = +2\hbar k\rangle \rightarrow |F = 2, p = +4\hbar k\rangle$. Reversing the effective wave vector of the beam splitters for the second pulse pair accomplishes both of these couplings. In principle, they are distinguished by a Doppler shift of $8\omega_r$ due to the speed difference between the lower and upper interferometer, as marked in Fig. I(c). Low-bandwidth beam-splitter pulses for atom interferometers typically resolve this frequency difference, but the high-bandwidth pulses we use to address a broad velocity class simultaneously address both transitions, unavoidably closing both interferometers. The two interferometer’s outputs ports (e.g. $B^-$, $B^+$) overlap spatially since the samples thermally expand faster than the interferometers separate.

We recover the recoil signal by using Raman beam splitters, which allow us to use state-dependent detection of the sum of signals from the lower and upper interferometers. Beginning in the $|F = 2\rangle$ ground state (state $A$) prior to the interferometry pulse sequence, the probability for an atom to emerge from the interferometer in the $|F = 1\rangle$ ground state (state $B$) oscillates as:

$$P_B = D [1 - C_+ \cos(\Delta \phi^-) - C_- \cos(\Delta \phi^+)], \quad (2)$$

where $C_\pm$ are the fringe contrasts of each interferometer and $D$ is an overall offset. For approximately equal contrasts, $C_+ = C_- \equiv C/2$, the signal simplifies to:

$$P_B = D[1 - C \cos(2ka_z(T + T') + 2\delta T) \cos(8\omega_r T)]. \quad (3)$$

Our setup is similar to the one previously described in Ref.\cite{38} but without the polarization gradient lattice used for sub-Doppler cooling. We heat lithium to 400°C and trap the vapor in a two-dimensional (2D) magneto-optical trap (MOT). A push beam tuned near resonance sends the atoms through a differential pumping tube into the interferometry chamber, where approximately 15 million atoms are trapped in a three-dimensional (3D) MOT. After lowering the intensities of both the cooling and repumping light and moving the detuning closer to resonance, the cloud reaches a final temperature of roughly 300 μK.

To define a quantization axis for optical pumping and Raman transitions, we apply a 1-G bias magnetic field along the $z$ axis. Despite the 250-μs decay of the current in the anti-Helmholtz MOT coils, the quadrupole field remains appreciable for milliseconds due to eddy currents in the steel vacuum chamber. We use the 3D MOT beams as optical molasses to limit the thermal expansion of the cloud while the eddy currents decay. No polarization gradient cooling occurs during this step due to the small detuning of the 3D MOT beams from the unresolved $D_2$ line (2$P_{3/2}$ state)\cite{38}.

After the optical molasses, the atoms are distributed among the five non-degenerate magnetic sublevels of the $|F = 2\rangle$ ground-state manifold. This leads to magnetic dephasing since the Ramsey-Bordé interferometer phase depends on the internal energies through the $\delta$ term. Interferometer experiments often select atoms in the desired magnetic sublevel by transferring them to the other hyperfine state with a microwave and blowing away the remaining populations with resonant light. The unresolved $D_2$ line in lithium, however, precludes the efficient cycling transitions required to impart the large momentum needed for such blow-away beams. Furthermore, this selection process is lossy, as large atomic populations are sacrificed to the blow-away beam.

To avoid the magnetic dephasing from atoms in dif-

![FIG. 2. Optical pumping: (a) Interference fringes without optical pumping (lower dashed curve) and with optical pumping (upper solid curve). Each gray point on the traces is the average of 5 experimental shots and error bars are omitted for clarity. (b) Optical pumping on lithium’s $D_1$ line with $\pi$ light (green arrow) results in a dark state at $|F = 2, m_F = 0\rangle$ (black circle). Atoms that decay to $|F = 1\rangle$ are recovered by 3D MOT repump light (yellow arrow). Each dash represents a unique magnetic sublevel.](https://example.com/figure2.png)
ferent magnetic sublevels, we optically pump the sample to the magnetically insensitive $|F = 2, m_F = 0\rangle$ state by taking advantage of the selection rule that prohibits $m_F = m'_F = 0$ transitions when $\Delta F = 0$. Once the magnetic field gradient decays below 1 G/cm (after 1.5 ms of optical molasses), we send 3 mW of light tuned within a linewidth ($\Gamma/2\pi = 5.87$ MHz) of the $|F = 2\rangle$ to $|F' = 2\rangle$ transition on the well-resolved $D_1$ line ($2P_{1/2}$ state). The optical pumping light is $\pi$ polarized along $\hat{z}$ and has a 3.6-mm Gaussian waist. Unlike the $D_2$ line, lithium’s $D_1$ line has a resolved hyperfine structure (see Fig. 2(b)). Optical pumping on the $D_1$ line therefore avoids the slightly off-resonant transitions ubiquitous on the $D_2$ line 39. In each of the six 3D MOT beams, we use 1.5 mW of $D_2$ MOT repump light to recover atoms that decay to $|F = 1\rangle$. We tune the repump frequency closer to resonance, optimizing for optical pumping efficiency. After 50 $\mu$s of optical pumping, more than 80% of the atoms occupy the dark state.

Figure 2 displays the efficacy of the optical pumping for interferometry. Without optical pumping, the recoil fringes have low contrast, a low signal-to-noise ratio, and decohere more rapidly, limiting the maximum interrogation time and sensitivity. Preparation to the magnetically insensitive $|F = 2, m_F = 0\rangle$ state before interferometry increases the contrast and signal-to-noise ratio by more than a factor of 2 at short interrogation times. Optical pumping also makes the fringes visible at longer interrogation times.

After optical pumping, we measure the fringes by varying the separation time $T$ while keeping $T' = 10 \mu$s and $\delta$ constant but small compared to $\omega_r$. To close the interferometers, we reverse the direction of the Raman beams for the second pulse pair using an electro-optic modulator (see Supplemental Material). For normalized detection, we use a new imaging technique that captures two images during a single exposure (see Supplemental Material).

Figure 3 shows the summed interference fringes obtained from the simultaneous conjugate Ramsey Bordé interferometers. As seen in Eq. 3 they can be described by a fast oscillation at a frequency of $8\omega_r$ within an envelope function that oscillates slowly at a frequency set by the two-photon detuning $2\delta$, in addition to accelerations of the atoms $a_z$. Here, the two-photon detuning term dominates over phases induced by acceleration, because we operate our interferometer perpendicular to gravity and at short interrogation times. Fig. 3(b) shows the fast component of the summed fringes. We fit the fringes using a least-squares method to the functional form in Fig. 3(b). The confidence interval in the fit constitutes a 32 ppm recoil measurement in 2 hours. After averaging across 10 such data sets with varying $\delta$, we reached a precision of 10 ppm. The phase sensitivity of the fit corresponds to a sensitivity roughly 50 times larger than the shot-noise limit.

The noise observed in the data is due mostly to laser noise, as we have confirmed by numerical simulations adapted from previous studies of noise in Ramsey-Bordé interferometers 41. The linewidth of the Raman laser ($\gamma/2\pi \approx 1$ MHz) is sizable compared to the small magnitude of the single-photon detuning ($\Delta/2\pi = 210$ MHz) and creates pulse-to-pulse fluctuations of the two-photon Rabi frequency, which result in noise significantly larger than the shot-noise-limited sensitivity.

The coherence time of the interferometer is not yet limited by thermal expansion out of the Raman beam but instead by magnetic dephasing of the $m_F = 0$ atoms. The magnetic field gradient that survives after the optical molasses gives rise to inhomogeneous quadratic Zeeman shifts, leading to an interferometer phase dependent on an atom’s position in the cloud. We are able to reduce the magnetic gradient by extending the optical molasses time to 5 ms and, with half the remaining gradient, the interference contrast indeed decays at half the rate. Magnetic gradient compensation would lead to longer coherence times and improved sensitivity. At a conservatively projected $T = 1$ ms, we estimate the shot-noise-limited sensitivity with $10^7$ atoms to be 100 ppb/√Hz. Implementing sub-Doppler cooling techniques 38, 41 to reach a temperature of 40 $\mu$K (approximately $8T_r$) would improve the sensitivity by $\sqrt{50}/8 \sim 3$, but still require the techniques in this paper.

Phase shifts due to vibrations cancel when the fringes
are summed in our detection scheme, as they enter the conjugate interferometers with opposite sign. The only effect of vibrations is then an amplitude modulation of the fringes. Consider Eq. (1) with a stochastic, Gaussian-distributed \( a_z \) with 0 mean and standard deviation \( \sigma \). When \( 2k\sigma T(T + T') \ll \pi \), the effect of such vibrations is a modulation of the interference contrast, which decreases proportionally to \( a_z^2 \). Other interferometers operating on a similar optical table without vibration isolation accrue phase shifts much less than \( \pi \) due to vibrations, even at \( T = 10 \text{ ms} \). Lithium’s high recoil frequency allows us to take sensitive data at \( T < 10 \text{ ms} \), and therefore to make full use of the common-mode rejection of vibration-induced signals.

This demonstration of interferometry opens the door to recoil measurements with other particles that are difficult to cool to subrecoil temperatures, such as electrons. Electrons, whose recoil frequency is on the order of GHz, are susceptible to relativistic effects and consequently a recoil-sensitive measurement can be used to measure Lorentz contraction [32]. While Kapitza-Dirac scattering has been proposed to realize matter-wave beam splitters for electrons in a Ramsey-Bordé interferometer [43], any vibrations or nonzero two-photon detuning will modify the phase \( \Delta \phi^- \) for a single Ramsey-Bordé. As we have shown in this work, the inclusion of the simultaneous conjugate interferometer (\( \Delta \phi^+ \)) recovers the recoil phase independently of a two-photon detuning even when the outputs of conjugate interferometers are spatially unresolved, as would be the case for electron plasmas in a Penning-Malmberg trap [32]. The required spectral resolution for detection could be achieved with bichromatic Kapitza-Dirac pulses. Bichromatic pulses with very large intensity have been proposed to impart momentum to an electron while inducing a spin flip [44] and hence couple the electron’s external and internal degrees of freedom. With such beam splitters acting on a spin-polarized sample and spin-dependent detection, the techniques we demonstrate in this work pave the way for a recoil-sensitive electron interferometer.

In summary, we demonstrate recoil-sensitive Ramsey-Bordé interferometry with laser-cooled lithium-7 at 300 \( \mu \text{K} \) (50Tc). The large Doppler spread of the sample is addressed with fast pulses, driving simultaneous conjugate interferometers with nearly equal contrast. Even with non-zero two-photon detuning, the interference fringes allow for the determination of the recoil frequency independent of two-photon detuning and vibrations. We suppress first-order magnetic dephasing and extend the coherence time by optically pumping the atoms to the magnetically insensitive \( |F = 2, m_F = 0 \rangle \) state using lithium-7’s well-resolved \( D_1 \) line. Our results relax cooling requirements for recoil interferometry, allowing for increased precision through high experimental repetition rates [31, 45]. Extending these techniques would allow for recoil-sensitive interferometry with atoms and other particles that have thus far been excluded from such experiments.

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We call this term a two-photon detuning loosely. It is the difference between the Raman laser frequency difference $\omega_1 - \omega_2$ and the hyperfine splitting $\omega_A - \omega_B$ during free evolution (see Fig. 1(b)). That is, our so-called two-photon detuning does not include AC Stark shifts of the internal hyperfine energies induced when the Raman light dresses the atoms, but it does include other shifts that persist during free evolution such as Zeeman shifts.
SUPPLEMENTAL MATERIAL

Raman beams: Light enters the setup as shown in Fig. 4 red-detuned with a single-photon detuning of $\Delta = 2\pi \times 210$ MHz from the crossover peak of the two hyperfine transitions on the $D_2$ line. Two acousto-optic modulators (AOM's) operating near 400 MHz (IntraAction ATM-4001A1) generate a frequency difference close to the 800-MHz ground state hyperfine splitting, each shifting either up or down by $(\omega_A - \omega_B + \delta)/2$. Approximately 30 mW of $\omega_1$ and 15 mW of $\omega_2$ coincide at the 2-mm diameter cloud in beams of 2.1-mm Gaussian waist. The beams realize a lin-lin geometry, with one polarized along $\hat{x}$ and the other along $\hat{y}$. The relatively small single-photon detuning allows for a high two-photon Rabi frequency $\Omega_R \sim 2\pi \times 1.6$ MHz and a short pulse duration. We drive a $\pi$ pulse in 320 ns with $\sim 30\%$ efficiency ($\tau_{\pi/2}=160$ ns). These large-bandwidth pulses address a considerable fraction of the atoms, whose two-photon resonance conditions are Doppler-broadened from the thermal velocity spread.

For Ramsey-Bordé interferometry based on Raman transitions, we must switch the propagation directions of $\omega_1$ and $\omega_2$ between the second and third pulses in order to close the interferometer. To achieve this, we orthogonally polarize $\omega_1$ and $\omega_2$ and overlap them before passing them through an electro-optic modulator (EOM) that acts as a voltage-controlled wave plate. A polarizing beam splitter following the EOM separates the frequencies and directs the light to fibers that send the beams to the atoms from opposing directions. By switching the EOM voltage from 0 V to 215 V during $T'$, we rotate the polarizations of the frequencies by $90^\circ$ and consequently reverse the Raman wave vectors.

Normalized detection with a single exposure: Shot-to-shot fluctuations in the total atomic population lead to noise in the interference signal. This effect can be lessened by normalizing each shot of the experiment to the number of atoms trapped on each shot. Interferometers based on cesium or rubidium can rely on state-selective cycling transitions to push a single hyperfine state aside and image spatially resolved populations. Lithium’s unresolved hyperfine structure precludes efficient cycling transitions, so another solution is needed. The long readout time of most CCD cameras makes it impossible to take successive exposures of hot samples, since the sample will have diluted by the beginning of the second exposure due to thermal expansion.

We normalize our detection with two state-selective images during a single 190-µs exposure of a CCD camera (PCO pixelfly) as described in Fig. 5. Light locked to the $D_2$ crossover passes twice through one of two 200-MHz AOM’s (Crystal Tech AOMO 3200-125), each of which produces an orthogonal polarization. Both frequencies address the $|F=1\rangle$ state on the $D_2$ line. During the first 90 µs of the exposure, we illuminate the atoms with one beam to image only the population in the $|F=1\rangle$ state. After a 10-µs delay, we switch on the same frequency with orthogonal polarization for 90 µs and turn on the 3D MOT cooling light. The cooling light depumps all atoms from $|F=2\rangle$ to $|F=1\rangle$ thus allowing us to detect the sum of the two states’ populations. The second image forms on the other side of the CCD, due to its deflection at the Wollaston prism. We allow the atoms to disperse and take a second (background) exposure with the same pulse sequence to generate side-by-side absorption images of the $|F=1\rangle$ population and the entire sample. The ratio of the two absorption imaging signals gives $P_{F=1}$. The extinction ratio in each polarization after the Wollaston is $\sim 10,000:1$, resulting in negligible crosstalk between the images. With one of the imaging beams blocked, no signal remains detectable due to the other beam in the final image.

![FIG. 4. Raman beam setup: 400-MHz AOM's shift the frequencies of the Raman beams to match the hyperfine splitting. The EOM reverses the propagation direction of the Raman beams with frequencies $\omega_1$ and $\omega_2$ between the second and third pulses, during $T'$.](image1)

![FIG. 5. Wollaston-based normalized imaging setup: orthogonal polarizations of the same imaging frequency are controlled by separate AOM's. The orthogonal polarizations form distinct absorption images during a single exposure of a CCD camera. The imaging beam is shown in red, while gray indicates the shadow that comprises the absorption imaging signal.](image2)