Atom interferometry using Kapitza-Dirac scattering in a magnetic trap

R. E. Sapiro, R. Zhang, and G. Raithel
FOCUS Center and Department of Physics, University of Michigan, Ann Arbor, MI 48109
(Dated: February 25, 2009)

We demonstrate two atom interferometric schemes based on Kapitza-Dirac scattering in a magnetic trap. In the first method, two Kapitza-Dirac scattering pulses are applied with a small time delay between them. High contrast interference is observed both using a thermal cloud and a Bose-Einstein condensate (BEC). In the second method, two Kapitza-Dirac scattering pulses are applied to a BEC with a time separation sufficiently large that the interfering orders complete half an oscillation in the magnetic trap; this enables interferometry between spatially separated paths.

PACS numbers: 03.75.Dg,37.25.+k,37.10.Jk,67.85.Hj

One of the most promising new technologies suggested by modern atomic physics is atom interferometry. Light interferometers have been used for numerous purposes, including probing materials (for an early example, see Ref. [1]), navigation (for reviews see Refs. [2, 3]), and searching for gravitational waves (see Ref. [4] for theory, and Ref. [5] and other LIGO publications for experimental status). Despite their many uses, however, light interferometers are mostly insensitive to electromagnetic and gravitational fields. Atoms, meanwhile, are sensitive to all of these fields; thus, an atom interferometer can provide a far more sensitive probe of force fields [4], in addition to being able to measure nearly anything a light interferometer can measure. Additionally, if the two arms of an atom interferometer are separately addressable, one can measure atom-surface interactions [6], electric polarizability [7, 8], and atom neutrality [9], as well as realize novel nanolithography schemes [10]. In this paper we present the experimental realization of two atom interferometry schemes using Kapitza-Dirac scattering. The first scheme can be used on thermal or Bose-Einstein condensed (BEC) atoms to produce high-contrast interference patterns with many fringes. The second scheme is used on a BEC to achieve a highly sensitive, large-separation atom interferometer, in which the maximal distance between the BEC components is sufficient that the components can be addressed separately.

Our experimental apparatus for cooling and trapping $^{87}\text{Rb}$ atoms is described in Ref. [11]. The atoms are trapped in an approximately harmonic magnetic potential with a frequency of 45 Hz along the axis of the interferometric arms, and 15 Hz and 45 Hz along the other two axes. We use an optical lattice to create two Kapitza-Dirac scattering pulses for splitting and recombining the atomic wavefunction. The lattice is composed of a retro-reflected laser beam with wavelength $\lambda = 852$ nm and lattice depth of about 80 recoil energies (one recoil energy is equal to $\hbar^2/2m\lambda^2$, where $m$ is the mass of $^{87}\text{Rb}$).

When a Kapitza-Dirac scattering pulse is applied to a plane matter wave, the wavefunction is split into several orders with momenta $2n p_r + p_0$, where $n$ is the scattering order, $p_0$ is the momentum before the Kapitza-Dirac scattering pulse is applied, and $p_r$ is the recoil momentum due to a single photon from the lattice ($p = \hbar/\lambda = m \times 5.4 \text{ mm/s}$). In our experiment, the value of $p_0$ varies randomly from shot to shot over a range $|p_0| \lesssim 0.4 p_r$, due primarily to magnetic field noise. The amplitude of the scattering order $n$ is given by $J_n(V_0\delta t/\hbar)$, where $J_n$ is the $n$th-order Bessel function of the first kind, $V_0$ is the full lattice depth, and $\delta t$ is the duration of the Kapitza-Dirac scattering pulse [12].

A schematic of the first interferometry scheme is shown in Fig. 1 (a). We use two Kapitza-Dirac scattering pulses: one to initially split the atomic wavefunction into mul-

\textbf{FIG. 1:} (a) Schematic of the first interferometry method, represented in momentum space. This method can be used for both a thermal cloud and a BEC. (b) Schematic of the second interferometry method, represented in position space. This method can be used to realize a large-separation atom interferometer.
tiple components and the other to recombine some of those components. When we split and recombine the components, we want the resultant interference pattern to arise from the interference of only two components. If more than two components interfere, the pattern becomes more complicated, and thus less straightforward to analyze. We also want the interfering orders to have approximately equal amplitudes, in order to maximize the visibility of the interference pattern. We choose $\delta t = 2 \mu s$, corresponding to $V_0 \delta t / \hbar = 1.6$. This causes only a small portion of the wavefunction to be in orders higher than 1 (∼6% of the wavefunction probability is scattered into each of the ±2 orders), and approximately equal amounts in the -1, 0, and +1 orders. After the application of the first Kapitza-Dirac scattering pulse, a time interval $\Delta t$, in the range of hundreds of microseconds, is allowed to pass. During this time, the ±1 orders move slightly in the trap due to the momentum imparted to them. After $\Delta t$, a second Kapitza-Dirac scattering pulse is applied. The wavefunctions in each order $n$ are split again into several orders $n'$. After both Kapitza-Dirac scattering pulses, the momentum of a component of the wavefunction is given by $2(n+n')p_n + p_0$ (neglecting the momentum change due to the magnetic trap).

With the second pulse, we again want to coherently split the wavefunction components into the orders $n' = 0$ and $n' = ±1$ with approximately equal amplitudes, while avoiding orders $|n'| ≥ 2$. Thus, we also pick $\delta t = 2 \mu s$ for the second pulse. Immediately after the second Kapitza-Dirac scattering pulse, the magnetic trap is turned off and the atoms undergo 20 ms of time-of-flight (TOF).

During the TOF, the wavefunction components spread into each other and interfere if they have approximately the same momentum, i.e. if their values of $n + n'$ are equal. If they have significantly different momenta, i.e. their values of $n + n'$ are different, they do not interfere, but rather gain macroscopic separation during TOF.

The description in the previous paragraph applies to a single atom interfering with itself. However, the experiment can be performed on a thermal cloud cooled to a sufficiently low temperature. At very low temperatures, the atoms in the thermal cloud will have similar enough momenta that even though the resulting measurement is an incoherent sum over all of the atoms, the interference pattern can still be seen [13].

We perform this experiment on a thermal cloud consisting of $2.35 \times 10^5$ atoms evaporatively cooled to just above the BEC transition. The results for various delay times $\Delta t$ are shown in Fig. 2 (a). As can be seen, interference fringes are clearly visible up to $\Delta t = 400 \mu s$. At image positions corresponding to momenta $±2p_r$ in Fig. 2 (a), two wavefunction components interfere with each other, namely the components $(n, n') = (0, ±1)$ and $(±1, 0)$. As a result, at these image positions the interference patterns exhibit regular, sinusoidal interference fringes. The fringe contrast varies with $\Delta t$; at $\Delta t = 50 \mu s$ a fringe contrast of over 90% is observed, while at $\Delta t = 400 \mu s$ the fringe contrast is only 33%. At the center of the images in Fig. 2 (a), three wavefunction components interfere with each other, namely $(n, n') = (0,0), (1, -1)$, and $(-1, 1)$, resulting in more complicated interference patterns. The interference patterns near zero momentum are most visibly complex when the fringe period is largest, at small $\Delta t$.

To obtain the fringe period, $\lambda_f$, from our experimental data, we determine the average distance between adjacent peaks for a given $\Delta t$, using multiple images, and ignoring the center region where the fringe pattern is more complex. The resulting plot, $\lambda_f$ vs $1/\Delta t$, is shown in Fig. 2 (b). The error bars shown in the figure are based on estimated reading uncertainties, due to camera resolution, and the number of samples. Applying a linear fit, we find a slope of 9.2 μm·ms⁻¹, with a fit uncertainty of ±1.1 μm·ms⁻¹.

The fringe period can also be calculated analytically and using simulations. To find $\lambda_f$ analytically, we assume the wavefunction is initially in the ground state of the harmonic potential (frequency $\omega = 2\pi \times 45$ Hz in our experiment). The first Kapitza-Dirac scattering pulse splits the wavefunction into several components with momenta differing by multiples of $2p_r$. Since the harmonic trap is left on between the Kapitza-Dirac scattering pulses, these components propagate without dispersion and maintain a position uncertainty of $\sigma = \sqrt{\hbar/2m\omega}$. The expectation values of the positions and momenta of the components follow the classical equations of motion of the harmonic
oscillator. The second Kapitza-Dirac scattering pulse, applied at time $\Delta t$, then splits the evolved wavefunction components a second time. After the second Kapitza-Dirac scattering pulse the harmonic-oscillator potential is turned off, and the wavefunction components propagate according to the well-known equation for a free Gaussian wave packet,

$$
\psi_l(x,t) \propto \frac{\exp(ik_l x - iE_l t/\hbar) \exp(-\frac{(x-x_1-k_l t/m)^2}{4\sigma^2(1+i\hbar t/(2\sigma^2 m))})}{\sqrt{2\pi(1+i\hbar t/(2\sigma^2 m))}}.
$$

(1)

Here, $x_1$ and $k_l$ are the expectation values of the position and momentum of component $l$ immediately after the second Kapitza-Dirac scattering pulse, and $t$ is the time elapsed after turning off the harmonic potential. The interference pattern after time-of-flight is obtained by setting $t = t_{TOF}$ and summing over all wave packet components. Specifically, the periodicity $\lambda_T$ of the interference pattern in the overlap region of two components, $l_1$ and $l_2$, follows

$$
2\pi/\lambda_T = |\frac{\partial \psi_{l_1}}{\partial x} - \frac{\partial \psi_{l_2}}{\partial x}|
$$

(2)

where $\psi_{l_1}$ and $\psi_{l_2}$ are the phases of wave packet components $l_1$ and $l_2$. After extracting $\psi_{l_1}(x,t)$ and $\psi_{l_2}(x,t)$ from Eq. 1 and taking the spatial derivatives, Eq. 2 leads to

$$
\lambda_T = \frac{2\pi}{m} \left( \frac{4m^2\sigma^4 + h^2\Delta t_{TOF}^2}{-4mk_{rel}\sigma^4 + h\Delta t_{TOF}x_{rel}} \right)
$$

(3)

with $k_{rel} = k_{l_2} - k_{l_1}$ and $x_{rel} = x_{l_2} - x_{l_1}$. Maximum interference contrast occurs at the classical intersection time of the wave packets, $t_{class} = -mx_{rel}/(hk_{rel})$; as expected, the fringe period $\lambda_T$ at that time equals $\lambda_{class} = 2\pi/k_{rel}$. Since the wave packets spread considerably during time of flight, for values of $t_{TOF}$ different from $t_{class}$ there can still be high-contrast interference, with fringe periods $\lambda_T$ different from $\lambda_{class}$. To find $\lambda_T$ for our experiment, we may consider, for instance, the interference of the scattering orders $(n,n') = (0,1)$ with $(n,n') = (1,0)$. For these, $k_{rel} = 2k_l(\cos(\omega\Delta t) - 1)$ and $x_{rel} = (2p_l/m\Delta t)\sin(\omega\Delta t)$. It is easy to verify that under the conditions of Fig. 2 the terms in Eq. 3 that involve $\sigma$ are much smaller than the terms that involve the time of flight, $t_{TOF}$. Noting further that under the conditions of Fig. 2, $\omega\Delta t \ll 1$, and thus $x_{rel} \approx 2p_l\Delta t/m$, we arrive at

$$
\lambda_T \approx \frac{h\Delta t_{TOF}}{2p_l\Delta t}
$$

(4)

This is the same equation as given in Ref. [13] for the case of free evolution between pulses and in the absence of mean-field interaction. Thus, in this first interferometry scheme, where $\omega\Delta t \ll 1$, the presence of the harmonic trap between the Kapitza-Dirac scattering pulses has no significant effect on the fringe period. According to Eq. 4, in our system we expect $\lambda_T = \frac{1}{\Delta t} \times 8.5 \mu\text{m} \cdot \text{ms}$. This is in reasonable agreement with our experimental observations.

The modeling described so far has been restricted to the case where the initial wavefunction is in the ground state, whereas the experiments discussed so far have been performed on thermal clouds. A thermal cloud of atoms at temperature $T$ in a harmonic oscillator (frequency $\omega$) can be modeled as an ensemble of minimum-uncertainty Gaussian wave packets with average initial positions $x_l$ and momenta $p_l$ following thermal distributions, $\omega^2m(x_l^2) = \hbar^2(k_l^2)/m = k_B T$. Forming weighted averages of the interference patterns produced by such wavefunction ensembles we find that high-contrast interference patterns may be observed for temperatures up to hundreds of nK. We have qualitatively confirmed this finding in the experiment by varying the temperature $T$ of the cold-atom cloud. We control the temperature by stopping the rf-induced evaporative cooling at different frequencies above the critical frequency at which a BEC begins to form. We determine the resultant temperature of the thermal cloud, after evaporative cooling and adiabatic relaxation of the magnetic trap, using TOF analysis along the vertical direction of the shadow images. We experimentally observe interference for temperatures up to $T \sim 650 \text{ nK}$. Due to uncertainties in the size of the initial atomic cloud before TOF, the uncertainties of the temperature measurements are $\sim 50 \text{ nK}$ (which corresponds to a fairly high relative uncertainty for the lowest temperature thermal clouds). In Fig. 3 we compare experimental interference patterns obtained with thermal samples to simulated patterns (simulated for the temperatures indicated in the figure), and find some qualitative agreement. However, the experimental results at high temperatures show greater fringe contrasts than the simulations; we still see contrasts of 13% at $\sim 650 \text{ nK}$ in the experiment, but the simulations indicate that the inter-
ference pattern should be completely washed out at that temperature. This parallels the observation of Miller *et al.* in Ref. [13]. Their proposed explanation of velocity-selective Bragg diffraction does not fit our experiment, though, since Kapitza-Dirac scattering is not velocity selective.

We have also implemented the described interferometry scheme using BECs of 5-8 × 10^4 atoms. Since the spatial extent of the BEC components after time of flight is much smaller than that of cold thermal clouds, it is more challenging to observe interference in BECs than in cold thermal clouds. Fringes in BECs can only be observed if the fringe period, \( \lambda_f \), is less than the size of the BEC components after TOF, but still larger than the spatial resolution of the imaging system (which is \( \sim 7 \mu \text{m} \) in our experiment). In our system, these contrasting requirements mean that BEC interference fringes can only be observed over the range 300 \( \mu \text{s} \lesssim \Delta t \lesssim 800 \mu \text{s} \). In Fig. 2 (a) we see that in the case \( \Delta t = 400 \mu \text{s} \), where there is a mix of thermal cloud and a very small BEC, a single interference fringe is only slightly smaller than a BEC component. In Fig. 2 (c), we show BEC interference for \( \Delta t = 420 \mu \text{s} \); several interference fringes are visible in each BEC component. For \( \Delta t \lesssim 300 \mu \text{s} \), the interference fringes are too wide to be visible, but any asymmetry between the paths of the +1 and -1 components will still cause a phase shift in the interference. These phase shifts manifest as one component having a greater amplitude than its symmetric component in the TOF image. An example leading to asymmetry between the paths of the +1 and -1 components is the initial momentum of the wavefunction, \( p_0 \). In our system, variations in the magnetic trapping fields cause a shot to shot fluctuation of \( p_0 \) over a range \( |p_0| \lesssim 0.4 p_r \). The resulting asymmetry in the interference pattern can be quantified by comparing the number of atoms \( P_- \) and \( P_+ \) with momenta -2\( p_r \) and +2\( p_r \), respectively, in the TOF images. In the presence of a cold thermal cloud, \( p_0 \) can be measured by using the interference of the thermal cloud as a reference. A thermal cloud has a sufficiently wide range of initial momenta that its interference pattern is independent of small variations in the average momentum. Thus, the center of the thermal cloud interference pattern can be taken as \( p = 0 \). The initial momentum of the BEC can then be found by comparing the position of the BEC component with \( n + n' = 0 \) to the \( p = 0 \) position as indicated by the thermal cloud. A straightforward calculation shows that, in the limit of small \( \Delta t \) and weak Kapitza-Dirac scattering (\( V_0 \delta t/\hbar \ll 2 A \)), the asymmetry of \( P_- \) and \( P_+ \), defined as \( A = (P_+ - P_-)/(P_+ + P_-) \), should vary as

\[
A = -\frac{\sin(2\Delta t k_f^2/m)\sin(2\Delta t k_{l_1}k_0/m)}{1 + \cos(2\Delta t k_f^2/m)\cos(2\Delta t k_{l_1}k_0/m)}
\]

where \( k_{l_1} = 2\pi/\lambda \) and \( k_0 = p_0/\hbar \). Using simulations, we have found that for \( \delta t = 2 \mu \text{s} \) (the value used in the experiments) Eq. 5 approximately holds for lattice depths, 2\( V_0 \), of up to about 70 recoil energies; for deeper lattices the dependence of \( A \) on \( \Delta t \) and \( k_0 \) becomes more complicated due to higher-order scattering. Therefore, to model our experiment, where the lattice depth is about 80 recoil energies, we use a numerical simulation. In Fig. 4, we plot \( A \) as a function of \( p_0 \) for images containing both a BEC and a thermal cloud with \( \Delta t = 100\mu \text{s} \) (squares), the simulation using propagators (solid line), and the analytic calculation (Eq. 5; dashed line). Since the lattice depth is large, the analytic method does not accurately describe our experiment. The experimental data do agree well with the theoretical result based on propagators, showing that the experimental shot-to-shot variations of \( p_0 \) result in well-characterized changes of the interferometric quantity \( A \).

The second interferometry scheme demonstrated in this paper differs from the first in that \( \Delta t \) is more than an order of magnitude larger: \( \sim 8 \text{ ms} \). This allows the non-zero momentum components produced by the first Kapitza-Dirac scattering pulse to complete a half-oscillation in the magnetic trap and return approximately to their initial positions near the center of the trap by the time the second Kapitza-Dirac scattering pulse is applied. This procedure, a schematic of which is shown in Fig. 1 (b), is a realization of a Mach-Zehnder interferometer. In the present work, clear demonstrations of this interferometer have only been possible using BECs. Under our experimental conditions, the +1 and -1 BEC components (starting with momenta of about +2\( p_r \) and -2\( p_r \), respectively) are, at the farthest point, 60 \( \mu \text{m} \) apart, which is much larger than the BEC size. This is far enough apart that it would be possible to manipulate the atoms traveling in one arm of the interferometer.
In the large-separation interferometer, the fringe period of the lattice is 60 Hz. The contrast of these fringes varies with the number of atoms in the magnetic potential (for the case described above, demonstrate that our large-separation interferometer has achieved larger \( \Delta t \)). For example, Refs. [14] and [15], 8 ms is still one of the larger time separations.

The images shown in Fig. 5, taken as described above, demonstrate that our large-separation interferometer produces interference fringes after a half round-trip of the atoms in the magnetic potential (for the case depicted in Fig. 5, the frequency of the trap along the axis of the lattice is 60 Hz). The contrast of these fringes varies considerably from shot to shot, ranging from 20% to 75%. In the large-separation interferometer, the fringe period is no longer described by Eq. 4 because \( \omega \Delta t \) is not \( \ll 1 \). At \( \Delta t = 8.4 \) ms, very close to a half-oscillation in the magnetic trap, \( \lambda_f \) is larger than a BEC component, and interference is observed indirectly as a difference in atom number between the components emerging with \( \pm 2p_n \), as described above for the other interferometry scheme. As seen in the figure, the +1 component has a larger population than the -1 component. At 7.2 ms and 7.8 ms, \( \lambda_f \) is small enough that multiple fringes can be seen in each BEC component. The presence of multiple fringes theoretically allows for more sensitive measurement, in that small changes in the phase of the fringes can be observed, rather than the less precise measurement of the number of atoms in each component. In order to exploit this potentially higher sensitivity, however, it would be necessary for the imaging system to be able to resolve small phase changes, which is not the case in our present setup.

The mean field energy of the BEC in our system is \( \sim h \times 100 \text{ Hz} \), so in theory it should be possible to begin to see small effects from the mean field after a half oscillation. However, although the number of atoms in our BEC varies by a factor of about two, we see no correlation between atom number and fringe phase when we compare images with the same \( p_0 \). When we simulate the interferometer and vary the mean field, we find that for the range of BEC sizes in our experiment, we would expect the variation in the fringe position due to the mean field to be smaller than the resolution of our camera, e.g. \( \sim \) 5 \( \mu m \) for the case of \( \Delta t = 7.8 \text{ ms} \). The simulations show a systematic shift of \( \sim 10 \mu m \), however, between a BEC of the size of the ones used in our experiment and one with an order of magnitude fewer atoms.

While our spatial and temporal separations are not as large as those in the experiment of O. Garcia et al. (who hold the record for a BEC interferometer with the arms completely spatially separate [14]), our system provides some advantages over theirs. One advantage is that we can make \( \lambda_f \) small enough to see individual fringes by making \( \Delta t \) not exactly equal to half an oscillation period in the magnetic trap. As mentioned above, this allows for potentially more precise measurements. Another advantage is that we use the curvature of our magnetic trap as the “mirrors” at the end of each arm of the Mach-Zehnder interferometer. This is both easier to implement and more reproducible than using a Bragg-scattering pulse. A Bragg-scattering pulse can lose its efficiency if the BEC is not exactly equal to half an oscillation period in the magnetic trap. As mentioned above, this allows for potentially more precise measurements. Another advantage is that we use the curvature of our magnetic trap as the “mirrors” at the end of each arm of the Mach-Zehnder interferometer. This is both easier to implement and more reproducible than using a Bragg-scattering pulse. A Bragg-scattering pulse can lose its efficiency if the BEC is not exactly equal to half an oscillation period in the magnetic trap. As mentioned above, this allows for potentially more precise measurements. Another advantage is that we use the curvature of our magnetic trap as the “mirrors” at the end of each arm of the Mach-Zehnder interferometer. This is both easier to implement and more reproducible than using a Bragg-scattering pulse.

The mean field energy of the BEC in our system is \( \sim h \times 100 \text{ Hz} \), so in theory it should be possible to begin to see small effects from the mean field after a half oscillation. However, although the number of atoms in our BEC varies by a factor of about two, we see no correlation between atom number and fringe phase when we compare images with the same \( p_0 \). When we simulate the interferometer and vary the mean field, we find that for the range of BEC sizes in our experiment, we would expect the variation in the fringe position due to the mean field to be smaller than the resolution of our camera, e.g. \( \sim \) 5 \( \mu m \) for the case of \( \Delta t = 7.8 \text{ ms} \). The simulations show a systematic shift of \( \sim 10 \mu m \), however, between a BEC of the size of the ones used in our experiment and one with an order of magnitude fewer atoms.

While our spatial and temporal separations are not as large as those in the experiment of O. Garcia et al. (who hold the record for a BEC interferometer with the arms completely spatially separate [14]), our system provides some advantages over theirs. One advantage is that we can make \( \lambda_f \) small enough to see individual fringes by making \( \Delta t \) not exactly equal to half an oscillation period in the magnetic trap. As mentioned above, this allows for potentially more precise measurements. Another advantage is that we use the curvature of our magnetic trap as the “mirrors” at the end of each arm of the Mach-Zehnder interferometer. This is both easier to implement and more reproducible than using a Bragg-scattering pulse. A Bragg-scattering pulse can lose its efficiency if the BEC is not exactly equal to half an oscillation period in the magnetic trap. As mentioned above, this allows for potentially more precise measurements. Another advantage is that we use the curvature of our magnetic trap as the “mirrors” at the end of each arm of the Mach-Zehnder interferometer. This is both easier to implement and more reproducible than using a Bragg-scattering pulse.

Based on the periodicity of the propagator of the one-dimensional harmonic oscillator, for given \( \Delta t_0 \) and integer \( n \) one would expect to find the same shadow images for all values \( \Delta t = \Delta t_0 + n \pi/\omega \). Thus, based on observations of thermal-cloud interference at some short time delay, \( \Delta t_0 \lesssim 500 \mu s \), as in Fig. 2, one might expect to find similar interference patterns of thermal clouds at much longer delay times, \( \Delta t = \Delta t_0 + n \pi/\omega \). We observe that cold, thermal clouds do indeed produce some interference after a half round-trip, as demonstrated in Fig. 5 (b). Nonetheless, in all cases studied the thermal-cloud interference observed after a half round-trip has been found to be much less pronounced than after short time delays, \( \Delta t \lesssim 500 \mu s \). The observed fringe contrast is 35% along the top edge of the shadow image, while it drops to zero.
near the center line and below. There are multiple causes that can explain the discrepancy between this observation and the high-contrast fringes observed for a BEC after the same time delay. The motion of the atoms in the trap may not be exactly one-dimensional. If the lattice laser is not perfectly aligned with one of the principal axes of the trap, the components gain a small momentum along a transverse axis of the trap. In this case, the interference planes in 3-D space may tilt after a large $\Delta t$ due to differences in trap frequency along the different axes of the trap. This would cause the interference to average out in the shadow images, which are 2-D projections of the atom density. This is also the probable cause of the noticeable tilting of the fringes in Fig. 5 (b). Additionally, any anharmonicity of the trap potential would break the periodicity of the propagator, leading to a reduction of the coherence time of cold, thermal samples.

Finally, as an extension of the previous scheme, we demonstrate that interference can be seen when the $\pm1$ components are allowed to undergo a full oscillation. This is shown for a BEC in Fig. 6, where $\Delta t = 16.5 \text{ ms}$.

In this case, the fringes appear tilted; this is presumably because the lattice laser is not perfectly aligned with the axes of the trap, as explained above. The full-oscillation version of the large-separation interferometer has both advantages and disadvantages over the half-oscillation version presented above. Having both components go through all the same space over the course of their separation negates any phase difference a component might pick up due to asymmetries in the trap. Time-independent, spatially-dependent fields and linear accelerations would affect both components equally, since they each sweep out the same spatial range. Each component is separately addressable only if the duration of the perturbation under investigation is timed such that it affects one component when it reaches one end of the trap but is no longer present by the time the other component reaches that location. Furthermore, when the $\pm1$ components collide at the center of the trap after half an oscillation, atom loss and four-wave mixing can occur due to s-wave scattering.

In conclusion, we have presented two new, simple atom interferometry schemes using Kaptiza-Dirac scattering. In one scheme, we can produce large, high-contrast interference fringes in both a BEC and a thermal cloud. In the other scheme we can see visible interference fringes after the BEC components follow spatially separated paths. This separation is sufficient for the BEC components to be individually addressable.

We acknowledge the support of AFOSR (grant FA9550-07-1-0412) and FOCUS (NSF grant PHY-0114336).