Experimental study of the role of trap symmetry in an atom-chip interferometer above the Bose–Einstein condensation threshold

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Abstract
We report the experimental study of an atom-chip interferometer using ultracold rubidium 87 atoms above the Bose–Einstein condensation threshold. The observed dependence of the contrast decay time with temperature and with the degree of symmetry of the traps during the interferometer sequence is in good agreement with theoretical predictions published in Dupont-Nivet et al (2016 New J. Phys. 18 113012). These results pave the way for precision measurements with trapped thermal atoms.

1. Introduction

Atom interferometers [1, 2] have demonstrated excellent performance in measuring gravity [3–6], gravity gradients [7–9] and rotations [10–12], using atoms in ballistic flight. In spite of being less well developed, trapped atom interferometers, for example using atom chips, [13, 14], would render the interrogation time independent of the atom’s flight, permitting miniaturization and possibly longer measurement times. To date, atom–chip-based interferometers have been successfully demonstrated using Bose–Einstein condensates [15, 16], but are subject to dephasing mechanisms resulting from atom–atom interactions [17–20]. Recently, we proposed a trapped atom interferometer using an ensemble of cold atoms above the Bose–Einstein condensation threshold (referred to as thermal atoms in the following) [21, 22]. This proposal is reminiscent of optical white light interferometry because of the necessity of keeping the path length difference between the two arms of the interferometer smaller than the coherence length. Similarly, the contrast decay of an atom-chip interferometer using thermal atoms is related to the degree of asymmetry between the two arms [22]. Related contrast decay effects have been described in optical traps, for example in [23, 24], where the asymmetry results from state-dependent light shifts, and [25], where the asymmetry is induced by spatial separation along the axis of a Gaussian beam.

In our experiment, the two arms of the interferometer correspond to two different internal states trapped in magnetic potentials, and we are able to control the effect of asymmetry without spatially separating the paths. The asymmetry can be tuned by adjusting the bias field, which results in slightly different magnetic moments as described by the Breit–Rabi formula. Using this technique in combination with evaporative cooling, we are able to measure the contrast decay time in an internal state interferometer (a Ramsey interferometer [26, 27]) for different values of the temperature and asymmetry, and compare our results with the theoretical predictions from [22].

The focus of this paper is inhomogeneous dephasing, as manifested by the contrast decay of Ramsey fringes [24]. On the other hand, homogenous dephasing, caused for example by fluctuating magnetic fields, and probed by spin echo measurements [24, 28], is beyond the scope of this paper. We first give a brief review of the theoretical predictions in section 2. Then, we describe in section 3 our experimental protocol and results. We finally compare the results to a simple model in section 4. We also discuss the identical spin rotation effect (ISRE) that was previously observed in similar experiments [29–33].
2. Theoretical model

In this section, we briefly recall the simple model of [22] describing the influence of the asymmetry on the contrast decay time. We consider a Ramsey interferometer involving two internal states of the $^{87}$Rb ground state manifold, namely $|F = 1, m_F = -1\rangle \equiv |a\rangle$ and $|F = 2, m_F = 1\rangle \equiv |b\rangle$, coupled by a two photon transition. During the whole interferometer sequence both states are maintained trapped, but not necessarily in identical potentials [16, 21], leading to inhomogeneous dephasing. We suppose that the traps are harmonic but with slightly different frequencies along one of the trapping axes, namely $\omega_a = \omega$ for state $|a\rangle$ and $\omega_b = \omega + \delta \omega$ for state $|b\rangle$ with $|\delta \omega| \ll \omega$. We also assume that the gas is at sufficiently high temperature $T$ to be accurately described by a Boltzmann distribution. The relative asymmetry $|\delta \omega|/\omega$ then implies an upper bound on the contrast decay time $t_c$, given by [22]:

$$t_c = \frac{\omega}{|\delta \omega|} \frac{\hbar}{kT}.$$  

(1)

Two differences between the experiment considered here and the model of [22] should be pointed out. First, in the model of [22], the relative asymmetry was assumed to grow linearly from zero to some finite value $|\delta \omega|/\omega$ where it was held for some interrogation time, before being ramped back to zero. It was furthermore assumed that the ramp was slow enough that the initial population of the eigenstates was conserved throughout the sequence. By contrast, the splitting and recombination are very fast in our experiment (on the order of the $\pi/2$ pulse time). Still, we expect the populations to be approximately conserved to a numerical factor on the order of unity as long as the relative asymmetry is small enough (typically smaller than $\hbar \omega/kT \gtrsim 10^{-3}$ in our experiment). Second, the model of [21, 22] is one-dimensional, while in the experiment described here the asymmetry occurs along all three trapping axes (with identical relative asymmetry, as will be seen below). In the three-dimensional harmonic case it can be shown that the contrast $C$ follows [34]:

$$C\left(\frac{t}{t_c}\right) = \frac{1}{\sqrt{(u^2 + 1)^3}},$$  

(2)

where $t$ is the interrogation time and $t_c$ is given by equation (1). This last equation could be extended to the case of Bose–Einstein distribution instead of a Boltzmann one. For the numbers considered in the following of this paper, at $t = t_c$ in the case of a Bose–Einstein distribution the contrast with differ from equation (2) by less than 4%.

3. Experiment

3.1. Tuning the asymmetry

The two interferometer states $|a\rangle \equiv |F = 1, m_F = -1\rangle$ and $|b\rangle \equiv |F = 2, m_F = 1\rangle$ are trapped by the same DC magnetic field. Because of the coupling between the nuclear angular momentum and the magnetic field, the effective magnetic moments (defined as the partial derivative of the energy with respect to the magnetic field) of the two states are slightly different, as described by the Breit–Rabi formula [35]. As pointed out in [36], there is a ‘magic’ magnetic field $B_m^0$ for which the effective magnetic moments of the two states $|a\rangle$ and $|b\rangle$ are identical. By changing the value of the minimum of the magnetic field around this value, we can go from a situation where the two traps are almost perfectly symmetric to a situation where they have significantly different frequencies.

Let us consider a static magnetic trapping field of the form $B(x) = B_0 + (\partial^2 B/\partial x^2)x^2$ where $\partial^2 B/\partial x^2 > 0$ and $x$ is a linear coordinate along one trapping axis. One can show [37], expanding the Breit–Rabi formula [35] up to the second order in the magnetic field, that the resulting trap frequencies are:

$$\omega_a \simeq \frac{\mu_B g_I \partial^2 B}{2m} \frac{\partial^2 B}{\partial x^2} \sqrt{1 + \frac{g_I}{g_J} \left(5 + \frac{4B_0}{B_m^0}\right)}$$  

(3)

for the state $|a\rangle$, and:

$$\omega_b \simeq \frac{\mu_B g_I \partial^2 B}{2m} \frac{\partial^2 B}{\partial x^2} \sqrt{1 + \frac{g_I}{g_J} \left(3 - \frac{4B_0}{B_m^0}\right)}$$  

(4)

for the state $|b\rangle$, where the magic magnetic field is given by [36]:

$$B_m^0 \simeq - \frac{8g_I E_{bd}}{3\mu_B(g_J - g_I)^2} \simeq 3.23 \text{ G}.$$  

(5)
In the above formulas, \( m \) is the atomic mass, \( \mu_B \) is the Bohr magneton, \( g_I \approx 2.002 \) and \( g_f \approx -9.95 \times 10^{-4} \) are the electron and nuclear spin \( g \) factors [35] respectively, and \( E_{\text{hfs}} \) is the energy splitting between the two hyperfine ground states of \(^{87}\text{Rb}\). The different curvatures for the two traps result in a non-zero value for the asymmetry along each trapping axis:

\[
\frac{\delta \omega}{\omega} = \frac{2|\omega_{A} - \omega_{B}|}{|\omega_{A} + \omega_{B}|} \approx \frac{4}{1 - g_I/(2g_f)} \left| \frac{g_I}{g_f} \left( \frac{B_0}{B_m} \right) \right|.
\]  

As expected, the asymmetry depends on the magnetic field \( B_0 \) at the trap minimum and vanishes when the latter is equal to the magic field \( (B_0 = B_m) \). The model can be easily generalized to 3 dimensions: the relative asymmetry is the same for all three trapping axes as can be seen in equation (6).

A simulation of our magnetic field geometry (performed using the known wire geometry on our atom-chip drawn in figure 1(a)) allows us to compute the potentials for the two states \( |A\rangle \) and \( |B\rangle \). Using a parabolic fit of the potentials we deduce the asymmetry. In figure 1(b), the asymmetry is shown as a function of the value of the magnetic field at the trap minimum, the model of equation (6) is in good agreement with the simulation. On the chip, a dimple trap is created by two crossing wires and an external bias field. We change the trap field minimum by changing the magnitude and direction of this bias field. In the whole span of \( B_0 \) in figure 1(b), the trap position changes by less than 1 \( \mu\text{m} \) and the trap frequencies, which are \( \sim85 \) and \( \sim148 \text{ Hz} \) in the horizontal plane and \( \sim161 \text{ Hz} \) in the vertical plane, by less than 5 Hz (numbers are given for state \( |B\rangle \)).

### 3.2. Experimental protocol and data

We use the experimental setup described in [37, 38] to trap a cloud of a few tens of thousands of \(^{87}\text{Rb}\) atoms in state \( |F = 2, m_F = 2\rangle \) on an atom-chip, using forced radio-frequency evaporation to control the final temperature, which is typically chosen between 150 and 800 nK (with the Bose–Einstein condensation threshold around 110 nK for state \( |b\rangle \)). After evaporative cooling in state \( |2, 2\rangle \), the cloud is transferred to state \( |b\rangle \) by microwave-stimulated Raman adiabatic passage as described in [38]. We then perform Ramsey spectroscopy between \( |a\rangle \) and \( |b\rangle \) for different values of the temperature and the asymmetry.

In figure 2, we show data for the longest \( t_e \) we have observed for \( B_0 \approx 3.247 \text{ G} \). In this case, given the cloud temperature of 200 nK equation (1) predicts \( t_e \approx 3.6 \text{ s} \left( \frac{\omega_e}{\omega} \right) = 10^{-5} \), the fitted value \( t_e = 2.6 \pm 1.9 \text{ s} \) is in good agreement. The interrogation time is limited to 600 ms so as not to overheat the atom-chip. For the magic field \( B_0 = 3.23 \text{ G} \), we observed a shorter \( t_e \) than for \( B_0 = 3.247 \text{ G} \) because the mean field corrections of the
resonant with the transition between the magnetic potentials, which empties the trap. To link our measurements and our simulation, we need to take into account the gravitational sag of the trap, thus we deduce the states change the magnetic field corrections compensate each other [33].

As described in the previous section, the bias field of the magnetic trap is tuned to investigate three different values of the asymmetry. For these three bias fields, the minimum of the magnetic field $B_0$ is measured (see table 1) by radio–frequency spectroscopy, i.e. we measure the frequency $\nu_i$, resonant with the transition between the states $|b\rangle$ and $|2, 0\rangle$ which empties the trap. To link our measurements and our simulation, we need to take into account the gravitational sag of the trap, thus we deduce $B_0$ by considering the energy at the minimum of the potential, i.e. at the position of the atoms: $E_{21}^{\text{fi}}(B_0) - mg^2/(2\omega^2) = \hbar\nu_iB_0$. Here $g$ is the standard acceleration due to gravity and $E_{21}^{\text{fi}}(B_0)$ is the energy difference (given by Breit–Rabi formula) between the states $|b\rangle$ and $|2, 0\rangle$ at the magnetic field minimum $B_0$. The linear potential caused by the acceleration of gravity does not change the magnetic field curvatures. Thus $B_0$ is used to calculate the asymmetry using the simulation shown in figure 1(b) (values are given in the column labeled ‘simulated asymmetry’ $\delta\omega/\omega$ in table 1).

For each of these three values of the asymmetry, the Ramsey signal is recorded for different values of the temperature. The local oscillator used for the excitation has a short term stability below $5 \times 10^{-12}$ at one second and long term drifts are corrected using a GPS clock signal. Examples of measured signals are shown in figure 3 for two different values of the temperature. The two outputs of the interferometer are measured via absorption imaging [37] and used to normalize the atom number in state $|b\rangle$ by the total atom number [39] (although only one is shown in figure 3). The highest value of $\delta\omega$ we achieved is one the order of 50 mHz, thus the first expected revivals [21] could occur at $1/\delta\omega \sim 20$ s, such Ramsey time can not be achieved in our experiment.

The envelopes of the fringes are extracted using a Hilbert transform [40] and then adjusted by a function of the form of equation (2) to infer the value of $t_i$ in equation (1) assuming asymmetry is the main source of contrast decay. The contrast decay time is plotted for the three values of the asymmetry as a function of the temperature in figure 4. Each of these curves is adjusted by a function of the form of equation (1) with $\delta\omega/\omega$ the fitting parameter. The resulting fitted values are given in table 1 in the column ‘inferred asymmetry’ $\delta\omega/\omega$.

### 4. Discussion

Our data is in reasonably good agreement with the model of equation (1), as can be seen in table 1. On the other hand, closer inspection of figure 4 indicates that the data points seem to decrease faster with the temperature than the adjusted law. This effect may point to trap anharmonicities which have been neglected up to this point. We can formulate a simple model of trap anharmonicity by considering the Hamiltonian:

$$\hat{H}_i = \frac{\hat{p}^2}{2m} + \frac{m\omega_i^2}{2} \hat{x}^2 + \frac{2}{\sqrt{15}} \sigma_i \hbar \omega_i \left( \frac{\hbar}{m\omega_i} \right)^{1/2} \hat{x}^3,$$

where $i$ labels the internal atomic state $|i\rangle$, $\sigma_i \ll 1$ is a dimensionless parameter quantifying the anharmonicity and $\hat{p}$ (respectively $\hat{x}$) is the momentum (respectively position) operator. The contrast decay time for such a

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**Table 1.** Trap asymmetry for different magnetic field minima $B_0$. We show the simulated values based on figure 1 and those inferred from the contrast decay time curves of figure 4.

| $B_0$       | Simulated asymmetry $\delta \omega/\omega$ | Inferred asymmetry $\delta \omega/\omega$ |
|-------------|------------------------------------|-------------------------------------|
| $2.471 \pm 0.020$ G | $(4.7 \pm 0.2) \times 10^{-4}$ | $(3.1 \pm 0.2) \times 10^{-4}$ |
| $2.874 \pm 0.019$ G  | $(2.2 \pm 0.3) \times 10^{-4}$ | $(6.7 \pm 0.9) \times 10^{-5}$ |
| $3.279 \pm 0.022$ G  | $(3.0 \pm 1.3) \times 10^{-5}$ | $(2.8 \pm 0.3) \times 10^{-5}$ |

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**Figure 2.** Ramsey fringes as function of the interrogation time $t_i$ for $B_0 = 3.247$ G. The fit of the envelope gives a contrast decay time $t_i = 2.6 \pm 1.9$ s.
potential is calculated in [22], leading to:

\[ \frac{1}{t^2} \sim 20(D\sigma)^2 \left( \frac{kT}{\hbar \omega} \right)^4 - 8D\sigma \delta \omega \left( \frac{kT}{\hbar \omega} \right)^3 + 3(\delta \omega)^2 + 2(D\sigma)^2 \left( \frac{kT}{\hbar \omega} \right)^2, \]  

(8)

where \( D\sigma = \omega \sigma^2 + \sigma^2 \delta \omega \) with \( \sigma^2 = (\sigma^2_1 + \sigma^2_2)/2 \) and \( D\sigma^2 = \sigma^2_1 - \sigma^2_2 \), \( \omega \) and \( \delta \omega \) are defined in section 2. Perfectly harmonic potentials correspond to \( D\sigma = 0 \) and we recover equation (1). A non vanishing \( D\sigma \) will tend to accelerate the contrast decay at high temperature. It was not possible to adjust the data to the above model because of an insufficient signal to noise ratio. However we can estimate an order of magnitude for the maximum value of the \( \sigma \) terms. For a temperature of 700 nK, \( kT/\hbar \omega \sim 120 \) (for \( \omega \) we took the geometric mean of the trapping frequencies), keeping only the first term, we find \( D\sigma/\omega = \sigma^2 + \sigma^2(\delta \omega/\omega) \lesssim (\hbar \omega)^2/(\hbar 20 \omega^2(kT)^2) \). In the most symmetric case (figure 4(c)) we observe a contrast decay time around 500 ms for temperature around 700 nK, thus \( D\sigma/\omega \lesssim 4 \times 10^{-8} \). This gives upper limits for the anharmonicity of the potentials: \( \sigma \lesssim 10^{-1} \), and for the difference between the cubic terms of the potentials: \( \delta \omega^2 \lesssim 4 \times 10^{-8} \). This is in agreement with a cubic fit of the simulation of our potentials which gives \( \sigma \sim 1 \times 10^{-2} \) and \( \delta \omega^2 \sim 10^{-5} \). 

It can be noticed from figure 3 that for some values of the asymmetry the phase coherence is lost even though the contrast remains appreciable (see for example figure 3(b)). We attribute this to magnetic field fluctuations.
Our experimental conditions (atomic density, temperature and inhomogeneity) are close to those of [32] where an ISRE results in long contrast decay times (∼60 s) and contrast revivals. One might wonder why this effect is not visible in our experiment. The ISRE regime is characterized by three time scales [32]: (i) the inhomogeneity $\Delta_0 / \omega_{ex} = kT \delta \omega / \hbar \omega \gamma \pi / 4$ with $\pi$ the mean atomic density and $\gamma = 4 \pi \hbar (a_{ss} - a_{aa}) / m$ where $a_{ss}$ is the scattering length between states $|i\rangle$ and $|j\rangle$, (ii) the elastic collision rate $\Gamma_{col} = (32 / 3) a_{ss}^2 \hbar \pi (\pi k T / m)^{1 / 2}$ and (iii) the exchange energy $\hbar \omega_{ex} = 4 \pi \hbar^2 a_{ss} \pi / m$. For ISRE to occur, the following three conditions must be satisfied [32]: (i) $\Delta_0 / \omega_{ex} \ll 1$, (ii) $2 \pi \Gamma_{col} / \omega_{ex} \ll 1$ and (iii) $2 \pi \Gamma_{col} / \omega_{min} \ll 1$ where $\omega_{min}$ is the smallest trapping frequency. To get a more quantitative criterion, we compare our experimental values of these parameters, for various densities, to the numbers from [32] where ISRE was observed.

In an attempt to reconcile our observations with those of [32], we developed a model whose details will be published elsewhere [34]. This model indicates that the trap anisotropy could significantly enhance ISRE. The ISRE arises from forward collisions between cold atoms which tend to be situated at the bottom of the trap and hot atoms which tend to sample magnetic fields far from the minimum. In the case of a cigar shaped trap, a hot

![Figure 4](image1.png)

Figure 4. Contrast decay time as a function of the atomic cloud temperature for three different values of the (inferred) asymmetry. The vertical scale (contrast decay time) is different in each plot. (a) $B_0 = 2.471$ G, corresponding to $\Delta_0 / \omega_{ex} = 1.5 \cdot 10^{-3}$, (b) $B_0 = 2.874$ G, corresponding to $\Delta_0 / \omega_{ex} = 2.2 \cdot 10^{-4}$ and (c) $B_0 = 3.279$ G, corresponding to $\Delta_0 / \omega_{ex} = 3.0 \cdot 10^{-5}$. The crosses stand for the experimental data with error bars, the solid blue lines are the fit of the data (inferred values and errors are given in table 1, column ‘inferred asymmetry’). The dashed black lines are the model with the simulated values of the asymmetry. The error bars on the temperature are one standard deviation of several repetitions of the measurement. The error bars on $t_c$ are a 66% confidence interval of the fitted value of $t_c$.

![Figure 5](image2.png)

Figure 5. The three relevant dimensionless parameters to assess the existence of ISRE. (a) $\Delta_0 / \omega_{ex}$, (b) $2 \pi \Gamma_{col} / \omega_{ex}$ and (c) $2 \pi \Gamma_{col} / \omega_{min}$ as a function of the mean atomic density $\bar{n}$ in unit of $10^{12}$ cm$^{-3}$. The three sets of crosses stand for our experimental data, while the open circles stand for the data of [32] where ISRE was observed.
atom oscillating in the trap comes close to the minimum of the trap at each oscillation enhancing forward collisions with cold atoms. In the case of an isotropic trap, a hot atom can oscillate without coming close to the trap minimum thus reducing the number of forward collisions with cold atoms and also the ISRE. Our trapping potential is closer to an isotropic trap (frequencies 85, 148, 161 Hz) than that in [32] (32, 97, 121 Hz), which might explain the discrepancy.

5. Conclusion

We have shown that the simple model derived in [22] describes reasonably well the contrast decay in our experiment, confirming the important role of symmetry in atom interferometry with thermal atoms. The next step is to introduce a spatial separation between the two internal states, for example by using near-field microwave gradients [16, 21, 22]. In this context, equation (1) could serve as a benchmark to assess the required degree of symmetry in the design of future atom-chip interferometers.

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