Sensitive gravity-gradiometry with atom interferometry: progress towards an improved determination of the gravitational constant

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\textit{New Journal of Physics} 12 (2010) 095009 (17pp)
Received 5 February 2010
Published 24 September 2010
Online at http://www.njp.org/
doi:10.1088/1367-2630/12/9/095009

Abstract. We present here the current status of our high-sensitivity gravity-gradiometer based on atom interferometry. In our apparatus, two clouds of laser-cooled rubidium atoms are launched in a fountain configuration and simultaneously interrogated by a Raman-pulse interferometry sequence. The system has recently been upgraded and its stability re-evaluated. We also discuss the recent progress of the experiment towards a precise determination of the Newtonian gravitational constant $G$. The signal-to-noise ratio and the long-term stability of the gravity gradiometer demonstrated interesting perspectives for pushing the $G$ measurement precision below the 100 ppm level.

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1. Introduction

Matter-wave interferometry has recently led to the development of new techniques for the measurement of inertial forces, finding important applications in both fundamental physics and applied research. The remarkable stability and accuracy that atom interferometers have achieved for acceleration measurements can play a crucial role in science and technology. Quantum sensors based on atom interferometry [1] underwent a rapid development during the last decade, and different schemes were demonstrated and implemented. Atom interferometry is used for precise measurements of gravity acceleration [2]–[5], the Earth’s gravity gradient [6, 7, 14] and rotations [8]–[11]. Currently, experiments based on atom interferometry are in progress to test Einstein’s Equivalence Principle [12, 13] and to measure the Newtonian gravitational constant $G$ [14]–[16]. In addition, experiments to test general relativity [13, 18] and the $1/r^2$ Newton’s law [17], [19]–[21], for a search of quantum gravity effects [22] and for gravitational wave detection ([23] and references therein, [24]) have been proposed. Accelerometers based on atom interferometry have been developed for many practical applications, including metrology, geodesy, geophysics, engineering prospecting and inertial navigation [7], [25]–[28]. The long interaction times possible in the space environment will allow us to take full advantage of the potential sensitivity of atom interferometers [29, 30].

Our atom interferometer MAGIA (Italian acronym for ‘Accurate Measurement of $G$ by Atom Interferometry’) was developed for a precise determination of the Newtonian gravitational constant $G$. The basic idea of the experiment and some preliminary results are presented in [14, 16, 31, 32].

The Newtonian gravitational constant $G$ plays a key role in the fields of gravitation, cosmology, geophysics and astrophysics and is still the least precisely known among the fundamental constants. Based on the weighted mean of eight values obtained in the past few years [33], in 2006 the Committee on Data for Science and Technology (CODATA) recommended a value with a relative uncertainty of 100 ppm. Although $G$ measurements have improved considerably since 1998 [34], the available values are still in poor agreement. Indeed, while the most precise measurements of $G$ have assigned uncertainties lower than 50 ppm [35]–[39], the results differ by many standard deviations. From this point of view,
the realization of conceptually different experiments can help us to identify still hidden systematic effects and therefore improve the confidence in the final result. With a few exceptions [38, 40, 41], most experiments were performed using conceptually similar schemes based on suspended macroscopic masses as probes and torsion balances or pendula as detectors. In our experiment, freely falling atoms act as probes of the gravitational field and an atom interferometry scheme is used to measure the effect of nearby well-characterized source masses. The projected accuracy for MAGIA shows that the results of the experiment will be important to discriminate between existing inconsistent values.

The paper is organized as follow. In section 2, we illustrate the principle of measurement; in section 3, we describe the apparatus, with special emphasis on recent upgrades; and in section 4, we present the experimental results.

2. Principle of measurement

A detailed description of the MAGIA experiment can be found in the previous papers [14, 16, 31, 42] and references therein. Here, we will recall the measurement principle and discuss the recent progress of the experiment.

In our experiment, we use atom interferometry to perform a simultaneous measurement of the differential acceleration experienced by two clouds of cold rubidium atoms in the presence of a well-characterized set of source masses. The measurement, performed for two different positions of the source masses, allows us to determine the Newtonian gravitational constant from the precise knowledge of the source masses distribution.

In a Raman interferometry-based gravimeter, atoms in an atomic fountain are illuminated by a sequence of light pulses that split, redirect and recombine the atomic wave packets. The light pulses are generated by two laser beams, whose frequencies $\omega_1$ and $\omega_2$ are resonant with the $\Lambda$-type transition of a three-level atomic system with two lower-energy states $|a\rangle$ and $|b\rangle$ and an excited state $|e\rangle$. The Raman laser beams, propagating along the vertical $z$-axis in opposite directions, are used to drive two-photon Raman transitions between $|a\rangle$ and $|b\rangle$.

Atoms are first prepared in the state $|a\rangle$. During the interferometer sequence, a $\pi/2$-pulse with duration $\tau = \pi/2\Omega$, $\Omega$ being the two-photon Rabi frequency, splits the atom wavefunction into an equal superposition of $|a\rangle$ and $|b\rangle$. The interaction with the Raman beams not only modifies the internal state of the atom but also results in a momentum exchange by an amount of $\hbar k_{eff} = \hbar (k_1 + k_2)$ ($k_i = \omega_i/c$; $i = 1, 2$) that modifies the atomic trajectories. After a time $T$, a $\pi$ pulse with a duration of $2\tau$ switches back the internal state from $|a\rangle$ to $|b\rangle$ and vice versa, redirecting the atomic trajectories. Finally, again after a time $T$, a $\pi/2$ pulse recombines the atomic packets in the two complementary output ports of the interferometer. At the output of the interferometer, the probability of detecting the atoms in the state $|a\rangle$ is given by $P_a = (1 - \cos \phi)/2$, where $\phi$ represents the phase difference accumulated by the wave packets along the two interferometer arms. In the presence of a gravity field, atoms experience a phase shift $\phi = k_{eff} g T^2$ depending on the local gravitational acceleration $g$ and on the time interval $T$ between the Raman pulses [2]. The gravity gradiometer consists of two absolute accelerometers operated in differential mode. Two spatially separated atomic clouds in free fall along the same vertical axis are simultaneously interrogated by the same Raman beams to provide a measurement of the differential acceleration induced by gravity on the two samples.
Figure 1. Scheme of the MAGIA experiment. $^{87}$Rb atoms, trapped and cooled in a MOT, are launched upwards in a vertical vacuum tube with a moving optical molasses scheme, producing an atomic fountain. Near the apogees of the atomic trajectories, a measurement of their vertical acceleration is performed by a Raman interferometry scheme. External source masses are positioned in two different configurations ($C_1$ and $C_2$) and the induced phase shift is measured as a function of mass positions.

3. Experimental apparatus

Figure 1 shows a schematic diagram of the MAGIA experiment. The gravity gradiometer setup and the configurations of the source masses ($C_1$ and $C_2$) are visible. At the bottom of the apparatus, a magneto-optical trap (MOT) with beams oriented in a 1–1–1 configuration collects $^{87}$Rb atoms. Using the moving molasses technique, the sample is launched vertically along the symmetry axis of the vacuum tube and cooled down to a temperature of about 2.5 $\mu$K. The gravity gradient is probed by two atomic clouds moving in free flight along the vertical axis of the apparatus and simultaneously reaching the apogees of their ballistic trajectories at 60 and 90 cm above the MOT. Such a geometry, requiring the preparation and launch of two samples with a large number of atoms in a time interval of about 100 ms, is achieved by juggling the atoms loaded in the MOT \[43\]. Shortly after launch, the two atomic samples are velocity selected and prepared in the ($F = 1, m_F = 0$) state using a combination of a Raman $\pi$ pulse and resonant blow-away laser pulses. The interferometers take place at the center of the vertical tube shown in figure 1. In this region, surrounded by two $\mu$-metal shields (76 dB attenuation factor of the magnetic field in the axial direction), a uniform magnetic field of 25 $\mu$T along the vertical direction defines the quantization axis. The field gradient along this axis is lower than 5 nT mm$^{-1}$. After the Raman interferometry sequence, the population of the ground state is measured in a chamber placed just above the MOT by selectively exciting the atoms on the $F = 1, 2$ hyperfine levels and detecting the resulting fluorescence.
Each atom interferometer in the gravity gradiometer measures the local acceleration with respect to the common reference frame identified by the wave fronts of the Raman lasers. Therefore, even if the phase noise induced by vibrations on the retroreflecting mirror completely washes out the atom interference fringes, the signals simultaneously detected on the upper and lower accelerometers remain coupled and preserve a fixed phase relation. As a consequence, when the trace of the upper accelerometer is plotted as a function of the lower one, experimental points distribute along an ellipse. The differential phase shift \( \Phi = \phi_u - \phi_l \), which is proportional to the gravity gradient, is then obtained from the eccentricity and rotation angle of the ellipse best fitting the experimental data [44].

The source masses are composed of 24 tungsten alloy (INERMET IT180) cylinders, for a total mass of about 516 kg. They are positioned on two titanium platforms and distributed in hexagonal symmetry around the vertical axis of the tube. Each cylinder is machined to a diameter of 100 mm and a height of 150 mm after a hot isostatic pressing treatment applied to compress the material and reduce density inhomogeneities. The two platforms can be precisely translated along the vertical direction by four step motors, with a resolution of 2 \( \mu \)m provided by an optical encoder; the positioning precision has been tested with a laser tracker [32].

In the preliminary measurement of \( G \) described in [16], the statistical uncertainty amounted to \( 1.6 \times 10^{-3} \), whereas the systematic uncertainty amounted to \( 4.6 \times 10^{-4} \). We have recently upgraded the MAGIA apparatus by addressing both the sensitivity and long-term stability to make them compatible with a measurement of \( G \) at the level of 100 ppm. In particular, we improved the intrinsic stability and long-term reliability of the whole laser system through a better frequency and intensity stabilization of the laser sources. We also implemented an atomic source based on a 2D-MOT (see section 3.1), and we installed a new Raman laser system (see section 3.2).

3.1. Two-dimensional magneto-optical trap (2D-MOT)

In the previous MAGIA setup, the magneto-optical trap was loaded from the background Rb vapor obtained from a dispenser. A major disadvantage of such an approach was the obvious trade-off between MOT loading rate and background pressure in the vacuum system. Indeed, a high Rb vapour density is important for a fast MOT loading, but it also degrades the vacuum, inducing higher atom losses during the ballistic flight and more background fluorescence at detection.

To achieve fast loading rates while preserving a very low background pressure in the MAGIA vacuum system, a high flux atomic source based on a 2D-MOT has been implemented [45]–[48]. Atoms evaporate from a temperature-controlled rubidium reservoir and, through a 15 mm diameter tube, they enter a \( 25 \times 25 \times 90 \) mm\(^3\) vapor cell (see figure 2). The cell is machined from a single piece of titanium and four rectangular windows (\( 15 \times 80 \times 3 \) mm) are glued to its sides for optical access. Two sets of coils are attached to the cell to provide radial magnetic gradients of about 0.2 mT mm\(^{-1}\). Two orthogonal beam pairs of cooling light with repumping light overlapped enter the vapor cell through the rectangular windows and radially cool the atoms. For the sake of compactness of the optical setup, each beam is split into three circularly polarized parts with a beam diameter of 24.5 mm. A low-intensity laser beam, slightly red detuned from the \( F = 2 \rightarrow F' = 3 \) cooling transition and propagating along the axial direction, pushes the atoms, increasing the flux in the direction of the 3D-MOT chamber. As a result, an atomic beam is coupled out through a hole (1.5 mm diameter) at the back wall of the
cell (2 mm thick). Before entering the UHV chamber, whose center is at about 0.5 m distance from the 2D-MOT, the generated atomic beam passes through a tube of purified graphite with a conical hole for differential pumping.

The laser system used to operate the 2D-MOT is based on a home-made Master Oscillator–Power Amplifier (MOPA) with an output power of about 500 mW. The master is an extended cavity diode laser using an interference filter for wavelength selection [49]. Two double-pass AOMs allow independent frequency and power tuning of the cooling and pushing beams. Optimal atomic flux from the 2D-MOT source is found when the optical frequencies of such beams are red tuned from the $F = 2 \rightarrow F' = 3$ transition by 8 and 13 MHz, respectively.

Under typical operating conditions, the measured values for the atomic beam flux, mean axial velocity, velocity spread and atomic beam divergence are $10^{10}$ atoms s$^{-1}$, 15 m s$^{-1}$, 7 m s$^{-1}$ and 23 mrad, respectively.

The 2D-MOT has reduced the background Rb density by more than two orders of magnitude in the UHV chamber. Figure 3 shows typical values of the MOT loading rate versus the temperature of the Rb oven in the 2D-MOT system. As compared to the standard operating
conditions of the dispenser previously employed to load the 3D-MOT, the MOT loading rate was increased by a factor of up to 5. The higher flux allows for a faster loading of the atomic clouds, with a consequent reduction in the measurement cycle, and/or allows for a higher number of detected atoms after the atom interferometry cycle, with a consequent increase in the signal-to-noise ratio (SNR).

3.2. The Raman laser system

The Raman beams are generated by two home-made MOPA systems with an output power of about 1 W each (see figure 4) [50]. Two interference filter stabilized extended cavity diode lasers are phase-locked with an offset frequency of about 6.8 GHz generated by a microwave synthesizer. In addition, one of the two lasers (master laser) is frequency-locked with an offset of about 2 GHz to the $F = 2 \rightarrow F' = 3$ transition, by detecting the beat note with a frequency stabilized laser (the MAGIA reference laser). Each laser beam injects an independent tapered amplifier (TA). The apparatus previously employed in the experiment was based on two grating tuned diode lasers in Littrow configuration, whose output beams were spatially overlapped to inject a single TA. The new system has several advantages. Indeed, interference stabilized extended cavity diode lasers have lower intrinsic frequency noise than Littrow grating stabilized lasers, improving considerably the locking time and stability. In addition, using two independent TAs instead of a single one allows independent manipulation and control of the two Raman beams. Such a scheme also provides higher optical power for the Raman beams. The higher Rabi frequency allows shorter $\pi/2$ and $\pi$ pulses, which should eventually lead to more atoms in the interferometer and higher contrast, i.e. a higher SNR. In our setup, the available power for Raman beams is about 240 mW.

Our Raman laser system features a double-stage optical phase-locked loop (OPLL). The primary OPLL detects the beat note between the two ECDL beams before injecting the TAs, in order to minimize the signal propagation delay, i.e. to maximize the loop bandwidth. The beat note is mixed with the 6.8 GHz reference frequency and the downconverted signal is

Figure 4. Scheme of the Raman laser system.
compared with a reference frequency generated by a direct digital synthesis (DDS) generator in a fast digital phase-frequency detector (Motorola MC100EP140). The DDS frequency is swept around 40 MHz with a phase-continuous linear frequency ramp to compensate for the change in Doppler effect during the interferometric sequence. The resulting error signal is filtered and used to drive two actuators on one of the ECDL (slave Raman laser), the voltage of the piezo holding the output coupler and the injection current of the laser diode. The loop bandwidth on the injection current is about 4 MHz. The output beams from the TAs are passed through two AOMs for independent power control and are finally recombined in a polarizing beam splitter. A third, single-pass AOM is used for pulse control just before coupling the Raman beams into an optical fibre for delivery to the atom interferometer.

We also apply an auxiliary, low-bandwidth loop in order to compensate for the phase noise introduced through the differential path of the two Raman laser beams before they are recombined in the optical fiber; to such purpose, we detect the beat note between the Raman laser beams at the polarizing beam splitter before AOM3; we mix the beat note with the 6.8 GHz reference frequency and we use another fast digital phase-frequency detector to compare the downconverted signal with the same ∼40 MHz reference frequency employed in the primary loop. The resulting error signal is properly filtered and used to control a voltage-controlled crystal oscillator (VCXO) driving the frequency of the AOM after the TA of the slave Raman beam. The resulting loop bandwidth is about 150 kHz. Figure 5 shows the phase noise spectral density measured in different conditions: after the primary loop, after the optical fiber with the secondary loop open and after the optical fiber with the secondary loop closed. The excess noise with respect to the primary loop is mainly due to the phase noise of AOM1 (see figure 4). The resulting phase noise is below 4 μrad/√Hz at frequencies above 10 Hz. With the typical parameters of our atom interferometry sequence (pulse duration τ = 24 μs, pulse separation T = 160 ms and repetition rate 0.56 Hz), the OPLL phase noise would limit the sensitivity of a single gravimeter to less than 10⁻⁹ g/√Hz [51]. Moreover, in a gradiometric measurement, the phase fluctuations of Raman lasers are largely suppressed as common mode noise.
4. Experimental results

In order to characterize the apparatus, we tested the sensitivity of the gravity gradiometer and its long-term stability. We also investigated the influence of Raman laser power and atomic flux on the instrument sensitivity.

Figure 6 shows a Lissajous figure obtained by plotting the normalized population of the $F = 1$ ground state detected at the output port of the upper interferometer as a function of the same measurement performed at the output port of the lower interferometer. The repetition rate of the experiment is about 0.55 Hz, and the plot contains 40 296 points.

We use a least squares fitting algorithm to extract the differential phase $\Phi_1$ of the gradiometer. We fit the experimental data to the parametric equations

$$
\begin{align*}
    x(t) &= A \sin(t) + B, \\
    y(t) &= C \sin(t + \Phi) + D,
\end{align*}
$$

where the $A$ and $C$ parameters represent the amplitudes of the interference fringes for the upper and lower interferometers, while $(B, D)$ are the coordinates of the ellipse center. This approach is adequate for the analysis of sensitivity and long-term stability. However, least square fitting yields a systematic error in the estimated value for $\Phi$ [52]; such a bias angle depends on both the noise and the ellipse angle itself; in particular, it is minimum for $\Phi \approx \pi/2$. More sophisticated algorithms have been proposed to retrieve $\Phi$ with Bayesian estimators. A Bayesian estimator has the advantage of being unbiased; however, it requires an a priori knowledge of the noise type. On the other hand, a least-square fit is affected by a bias for high noise level, but does not require a noise model. Although a discussion of systematic uncertainty in the $G$ measurement is beyond the scope of this paper, we have determined the corresponding bias as a function of the noise level with a simulation, and we have compared it with a Bayesian estimation assuming additive Gaussian noise; according to our results, in our experimental conditions the bias can be controlled at the level of a few tens of $\mu$rad.

Figure 6. Lissajous plot of 40 296 data points.
4.1. Influence of Raman laser power

We analyzed the effect of Raman laser power on the gradiometer sensitivity. For this purpose, we measured the differential interferometer phase $\Phi$ using two different values for the overall Raman laser power along the atom interferometry pulse sequence. Figure 7 shows two ellipses corresponding to Raman laser powers of 240 and 120 mW. Pulse widths were consequently adjusted to match the area condition for $\pi$ and $\pi/2$ pulses. The Raman laser power during the velocity selection pulse was kept at 120 mW in the two measurements, to ensure the same number and velocity distribution of the atoms prepared on the $F=1$ state before the 3-pulse interferometric sequence. The power ratio of Raman laser beams was chosen in such a way as to cancel the first-order light shift [53].

Each ellipse in figure 7 contains 960 points. Shorter Raman pulses during the interferometric sequence are indeed able to address a velocity class larger than that of the $F=1$ atoms after velocity selection. This improves the contrast of atom interference fringes and reduces the rms noise on the fitted ellipse angle $\Phi$ from 3 mrad at 120 mW to 2 mrad at 240 mW. Thus, doubling the Raman laser power improves the gradiometer sensitivity by about 30%.

4.2. Influence of atomic flux

We analyzed the effect of atomic flux on the gradiometer sensitivity. For this purpose, we measured the differential interferometer phase $\Phi$ using five different values for the loading rate from the 2D-MOT. For each value of the flux, we acquired 960 ellipse points. Figure 8 shows the rms noise on the fitted ellipse angle $\Phi$ versus the average number of detected atoms in the $F=2$ state. By fitting the experimental data with a power law, we find that the rms noise $\delta \Phi$ scales with the number of atoms $N$ as $\delta \Phi \propto N^{-0.54}$. We evaluate the number of
atoms from the number of detected photons, the overall detection efficiency, and the probe beams intensity and detuning. The expected scaling would be \(1/\sqrt{N}\) if we were limited by quantum projection noise (QPN). However, when detecting \(N = 7 \times 10^4\) atoms, the noise on atomic population measurement corresponds to about three times the QPN limit, indicating that other noise sources are dominant. Our measurements could be explained by assuming the contribution of two additional noise sources: a technical noise inversely proportional to \(N\) and an uncorrelated detection noise independent of \(N\). By fitting our data with the model

\[
\delta \Phi(N) = \sqrt{\left(\frac{A}{\sqrt{N}}\right)^2 + \left(\frac{B}{N}\right)^2 + C^2}
\]  

(2)

under the constraint \(\delta \Phi(N) \simeq 3 \times A/\sqrt{N}\), we obtain \(A = 0.2\ \text{rad} \sqrt{\text{atom}}, B = 80\ \text{rad atom}\) and \(C = 1.6\ \text{mrad}\). The relative contributions of the three noise terms in equation (2) are shown in figure 8.

4.3. Gradiometer sensitivity

As a test of the sensitivity of our apparatus, we observed the statistical fluctuations of the gradiometer measurements over about 20 h, keeping the masses in a fixed position. The data in figure 6 have been divided into a series of \(n\) consecutive data points. Each group of \(n\) points was fitted with an ellipse and the value for \(\Phi\) has been extracted with its estimated error. We have evaluated the Allan deviation of \(\Phi\) and verified that it scales as the inverse of the square root of the integration time, showing the typical behavior expected for white noise. Figure 9 shows the Allan plot for \(n = 24\).

The optimal number of points required to fit an ellipse has been estimated by varying the number of points \(n\) per ellipse and computing the Allan deviation for the series of \(\Phi\) values. As a figure of merit, we consider \(\eta(n) = \sigma(1 \text{ s})\), i.e. the Allan deviation at 1 s, calculated by splitting the data set into \(n\) points per ellipse. Since the fit is heavily nonlinear, \(\eta(n)\) drops sharply at first and then it reaches a plateau, as shown in figure 10. For large values of \(n\), \(\eta(n)\) seems to

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Figure 8. The rms noise on \(\Phi\) versus the number of detected atoms; each value is obtained from 960 data points; the solid line is a power-law fit, while the dashed line is a fit with equation (2), and the dash-dotted lines are the three corresponding terms of equation (2).
Figure 9. Allan deviation of $\Phi$ calculated from the data shown in figure 6, using 24 points per ellipse.

Figure 10. (Top) Allan deviation at 1 s of the ellipse angles calculated from the data shown in figure 6, versus the number of points per ellipse; (bottom) the average ellipse angle $\Phi$ calculated from the data shown in figure 6, versus the number of points per ellipse.

increase, although the error bars are quite large due to the small number of available points. This effect is not in contradiction with the absence of drifts in the $\Phi$ measurement over 20 h (see figure 9) and can be attributed to long-term drifts of the ellipse center. In fact, a similar plot for the average phase angle $\Phi(n)$ shows that as $n$ increases, the fit stabilizes to a well-defined phase value. By choosing $n$ as the smallest value that reaches the plateau, both S/N and temporal resolution are optimized. With $n = 200$, we have 201 ellipses from the data shown in figure 6. In such conditions, we measure a sensitivity of 90 mrad at 1 s, corresponding to a sensitivity to differential accelerations of $2 \times 10^{-8} g$ at 1 s. The resulting sensitivity is about a factor of 2.
Figure 11. (Top) Modulation of the differential phase shift measured by the atomic gravity gradiometer when the distribution of the source masses is alternated between configuration $C_1$ (upper points) and $C_2$ (lower points). Each point in the upper graph is the weighted average of five consecutive phase measurements obtained by fitting a 192-point scan of the atom interference fringes to an ellipse. (Bottom) Resulting values of the angle of rotation $\Delta\Phi(i)$; dashed horizontal lines mark the $\pm2\sigma$ confidence interval for the average.

better than in [16], thus reducing by a factor of 4 the integration time needed to reach a specific precision in the $G$ measurement. The regime of 100 ppm uncertainty can now be reached in about 65 days of continuous measurement, assuming that the error keeps averaging down as $1/\sqrt{\tau}$ on that time scale.

We also modulated the position of the source masses, as shown in figure 1. Figure 11 shows a measurement of the differential interferometric phase on a period of 120 h over 2 weeks. We moved the masses from the close ($C_1$) to the far ($C_2$) configuration and vice versa every 30 min, corresponding to 960 measurement cycles. On the time scale of the mass modulation period, we expect our gradiometer to be immune from drifts, as demonstrated by the measurement of figure 9. We split each set $C_n(i)$ of 960 points ($n = 1, 2$) into five series of 192 consecutive points; we fitted each group of five ellipses and we evaluated the average angle $\Phi_n(i)$ and the standard error on the average $\delta\Phi_n(i)$. In the sequence $C_1(i-1), C_2(i), C_1(i+1)$, we compared $\Phi_1(i)$ with the average of $\Phi_2(i-1)$ and $\Phi_2(i+1)$, to remove possible residual linear drifts in time. In this way, we obtained 119 couples of data $\{\Phi_1(i), \Phi_2(i)\}$. From each couple, a value for the rotation angle $\Delta\Phi(i)$ due to the position of the source masses can be obtained. The final result is $\Delta\Phi = 0.54597 \pm 0.00022$ rad with a $\chi^2$ of 128. This corresponds to a statistical uncertainty of $4.0 \times 10^{-4}$ on the measurement of $G$. The Allan deviation of the 119 values $\Delta\Phi(i)$, obtained by removing the time gaps, scales as the inverse root of the number of measurements (see figure 12), as expected under the assumption that the dominating noise on the measurement of $\Delta\Phi$ is white over the time scale of 2 weeks.

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Figure 12. Allan deviation of the $\Delta \Phi$ values shown in the lower plot of figure 11, after removal of the time gaps. The solid line is a fit with an inverse square root function of the number of measurements; each measurement lasts for 1 h.

5. Conclusions

We presented a sensitive gravity gradiometer based on Raman atom interferometry. Recent upgrades to the MAGIA apparatus have allowed us to reach a sensitivity to differential gravity accelerations of $2 \times 10^{-8} \text{ g s}^{-1}$.

We also discussed the system performance for a measurement of the Newtonian gravitational constant. The MAGIA experiment can run continuously for several days, showing a reproducibility of the gravity gradient measurement compatible with the stated sensitivity on the time scale of a few weeks. Our measurement of the differential gravity gradient can reach a statistical uncertainty of $4.0 \times 10^{-4}$ on the measurement of $G$ after 120 h of integration time.

A measurement of $G$ by atom interferometry at the level of 100 ppm is within reach. With the demonstrated sensitivity, an integration time of about 65 days would be required for reaching such uncertainty levels. With careful optimization of the current experimental configuration, we expect to reduce the required integration time below a couple of weeks. This can be attained by reducing the technical noise of detection, thus approaching the QPN limit, and by increasing the number of atoms through a cleaner preparation of the atomic state before the interferometric sequence. Further improvements in the setup, such as the use of more sophisticated detection schemes [54] or the implementation of high-momentum beam splitters [55], may enable an even higher sensitivity.

Acknowledgments

This work was supported by the INFN (MAGIA experiment), the EU (FINAQS STREP/NEST project contract no. 012986) and the ESA (SAI project contract no. 20578/07/NL/VJ). FS acknowledges financial support from the ESF through the Euroquasar program. GMT acknowledges useful suggestions from M Kasevich and A Peters. J Stuhler, T Petelski, A Bertoldi, G Lamporesi, M Jacquey, M de Angelis and A Giorgini contributed to setting up the experiment.
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