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## Controlling the Cold Collision Shift in High Precision Atomic Interferometry

F. Pereira Dos Santos, H. Marion, S. Bize, Y. Sortais, and A. Clairon

*BNM-SYRTE, Observatoire de Paris, 61 Avenue de l'Observatoire, 75014 Paris, France*

C. Salomon

*Laboratoire Kastler Brossel, ENS, 24 rue Lhomond, 75005 Paris, France*

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We present a new method based on a transfer of population by adiabatic passage that allows one to prepare cold atomic samples with a well-defined ratio of atomic density and atom number. This method is used to perform a measurement of the cold collision frequency shift in a laser cooled cesium clock at the percent level, which makes the evaluation of the cesium fountain accuracy at the  $10^{-16}$  level realistic. With improvements, the adiabatic passage would allow measurements of density-dependent phase shifts at the  $10^{-3}$  level in high precision experiments.

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Collisions play an important role in high precision atomic interferometry [1,2]. In most experiments, a precise control of the atomic density is hard to achieve, which sets a limit to how accurately systematic effects due to collisions can be corrected for. This is particularly true for clocks using laser cooled Cs atoms [3,4]. The accuracy of cesium fountains now reaches  $\sim 1 \times 10^{-15}$  [5]. It is currently limited by a 10%–20% systematic and statistical uncertainty on the determination of the cold collision shift [6–8]. To reach such an accuracy, one has to operate with a moderate number of detected atoms, typically  $10^5$ , which sets a standard quantum limit to the frequency stability [9] of about  $10^{-13} \tau^{-1/2}$ , where  $\tau$  is the averaging time in seconds. However, when using a high number of atoms ( $10^6$ – $10^7$ ), a stability approaching  $3 \times 10^{-14} \tau^{-1/2}$  has already been demonstrated [7,9], which would make the evaluation at the  $10^{-16}$  level practicable. Under these conditions, the cold collision frequency shift is very large ( $\sim 10^{-14}$ ). To actually reach such an accuracy, this shift has to be determined more precisely than presently achieved.

In this Letter, we present a method using adiabatic passage (AP) [10,11] that allows one to prepare atomic samples with well-defined density ratios. This enables the determination of the collisional frequency shift at the percent level or better, in an atomic clock or atom interferometer.

The measurement of the cold collision shift is based on a differential method [12]. One alternates sequences of measurements with two different, high and low, atomic densities. The frequency difference between the two situations, as well as the difference in the number of detected atoms, is measured. In a fountain, the collisional frequency shift is proportional to the effective density, as defined in [12]. Unfortunately, one has no direct access to this density. This density, and thus the collisional frequency shift  $\delta\nu_{\text{col}}$ , are still proportional to the number of detected atoms  $N_{\text{det}}$ :  $\delta\nu_{\text{col}} = KN_{\text{det}}$ . A linear extrapolation

of the clock frequency to zero detected atoms is valid only if the  $K$  coefficients for the high and low densities are the same.

Up to now, two methods have been used to change the density, and hence  $N_{\text{det}}$ . Atoms are initially loaded in an optical molasses, whose parameters (duration, laser intensity) can be varied. Another technique consists in keeping the same loading parameters but changing the power in a selection microwave cavity, which is used to prepare atoms in the  $|F = 3, m_F = 0\rangle$  state. One can select all (half) of the atoms initially in the  $|F = 4, m_F = 0\rangle$  state by applying a  $\pi$  ( $\pi/2$ ) pulse and pushing away the atoms remaining in the  $|F = 4\rangle$  state. However, because of the microwave field inhomogeneities in the cavity, the pulses cannot be perfectly  $\pi$  and  $\pi/2$  pulses for all the atoms. Besides affecting the atomic densities, both techniques also affect position and velocity distributions, as well as collisional energy, and consequently the  $K$  coefficients usually differ for the low and high density cases. Our numerical simulation, which takes into account the fountain geometry, the position, and the velocity distributions of the atomic cloud, shows that, for these methods, a linear extrapolation leads in the best case to a 10%–20% error in the determination of the collisional shift. Fluctuations and imperfect determination of those parameters prevent one from performing an accurate evaluation of the  $K$  coefficient. Moreover, for cesium atoms, as the collisional shift is energy dependent [13], this error can be even larger.

A method immune against these systematic effects prescribes to change the number of atoms of the sample without changing either its velocity distribution or its size. This can be realized by an adiabatic transfer of population, which allows one to prepare two atomic samples, where both the ratio of the effective densities and the ratio of the atom numbers are exactly 1/2. In contrast to previous methods, this one is insensitive to fluctuations of experimental parameters such as the size

and temperature of the atomic sample, or the power coupled into the selection cavity.

First, an adiabatic passage in the selection cavity is used to transfer with a 100% efficiency all the atoms from the  $|F = 4, m_F = 0\rangle$  state to the  $|F = 3, m_F = 0\rangle$  state [10,11]. This requires that the frequency of the microwave field in the cavity is swept across resonance, and that the Rabi frequency  $\Omega(t)/2\pi$  has an appropriate shape and maximum value,  $\Omega_{\max}/2\pi$ . We use Blackman pulses (BP), which minimize off-resonance excitation [14]. In order to fulfill the adiabaticity condition, the detuning from resonance  $\delta$  has to be chirped according to

$$\dot{\delta}(t) \propto \Omega^2(t). \quad (1)$$

Figure 1 shows the evolution of the microwave field amplitude together with the frequency chirp.

Second, we exploit another striking property of AP: if we stop the AP sequence when  $\delta = 0$  [half-Blackman pulse (HBP)], the atoms are left in a superposition of the  $|F = 4, m_F = 0\rangle$  and  $|F = 3, m_F = 0\rangle$  states, with weights rigorously equal and independent of the Rabi frequency. After removal of the  $|F = 4\rangle$  atoms with a pushing laser beam, half of the atoms are in the  $|F = 3, m_F = 0\rangle$  state, as desired.

In order to optimize this AP method and to evaluate its sensitivity to experimental parameters, we first performed a simple numerical simulation, solving the time-dependent Schrödinger equation for a two level atom in a homogeneous microwave field. The choice of the pulse parameters comes from a compromise between the insensitivity of the transition probabilities to fluctuations and the parasitic excitation of nonresonant transitions. Figure 2 displays as lines the calculated transition probabilities as a function of the maximum Rabi frequency  $\Omega_{\max}/2\pi$ , for BP and HBP. The parameters were a duration  $\tau_{\text{int}} = 4$  ms and  $\delta_{\max} = 5$  kHz, which were constant over the course of the experiment. The simulation shows that the transition probabilities deviate from 1

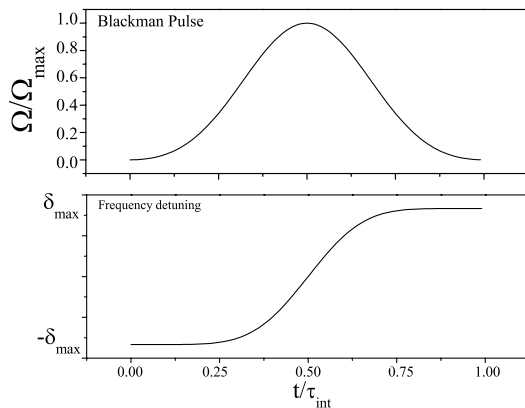


FIG. 1. Temporal dependence of the Blackman pulse and the corresponding frequency chirp.

and  $1/2$  by less than  $10^{-3}$  as soon as  $\Omega_{\max}/2\pi$  is larger than 2.4 kHz. A more realistic simulation has been performed, taking into account the Gaussian spatial distribution of the atomic cloud (characterized by  $\sigma = 3.5$  mm) and its trajectory, as well as the microwave field distribution of the  $\text{TE}_{011}$  mode for our cylindrical selection cavity. The simulation indicates that the transition probability deviates from the ideal value by less than  $10^{-3}$  for both BP and HBP for all the atoms contained within  $\pm 3.5\sigma$  of the vertical spatial distribution (more than 99.95% of the atoms in the cloud). The simulation indicates that in order to minimize the sensitivity of the transfer efficiency to microwave field inhomogeneities and timing errors, the center of the cloud has to be located in the center of the cavity at the end of the HBP. For instance, a delay as large as 1 ms with respect to the optimal timing induces only a  $7 \times 10^{-5}$  variation on the transition probability. The only critical parameter is the accuracy of the detuning at the end of the chirp for HBP. We calculate a linear sensitivity of the transition probability to the final detuning of  $7 \times 10^{-5} \text{ Hz}^{-1}$  for  $\Omega_{\max}/2\pi = 7.5$  kHz.

We use one of our Cs fountains to demonstrate the AP method and the resulting ability to control the collisional shift. This clock is an improved version of the Rb fountain already described elsewhere [7,15]. We use a laser slowed atomic beam to load about  $10^9$  atoms within 800 ms in a lin  $\perp$  lin optical molasses, with six laser beams tuned to the red of the  $F = 4 \rightarrow F' = 5$  transition at 852 nm. The atoms are then launched upwards at  $\sim 4.1$  m/s within 2 ms, and cooled down to an effective temperature of  $\sim 1 \mu\text{K}$ . After launch, the atoms are prepared into the  $|F = 3, m_F = 0\rangle$  state using a combination of microwave and laser pulses: they first enter the selection cavity ( $Q \sim 1000$ ) tuned to the  $|F = 4, m_F = 0\rangle \rightarrow |F = 3, m_F = 0\rangle$  transition, where they experience either BP or

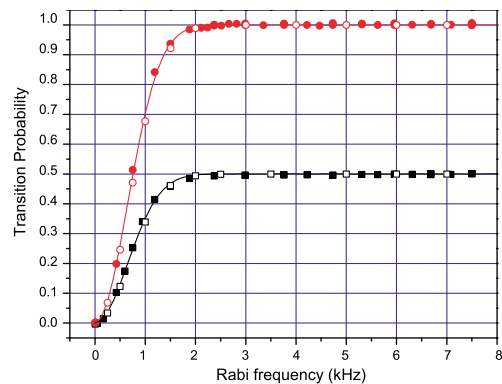


FIG. 2 (color online). Transition probabilities for a BP and a HBP with  $\tau_{\text{int}} = 4$  ms,  $\delta_{\max} = 5$  kHz as a function of  $\Omega_{\max}/2\pi$ . The results of the numerical simulations are displayed as lines (homogeneous Rabi frequency case) and open symbols ( $\text{TE}_{011}$  cavity case), whereas the measurements are full symbols.

HBP pulses. The atoms left in the  $F = 4$  state are pushed away by a laser beam tuned to the  $F = 4 \rightarrow F' = 5$  transition, 10 cm above the selection cavity. The amplitude of the Blackman pulses are shaped by applying an adequate voltage sequence (500 steps) to a microwave voltage-controlled attenuator (60 dB dynamic range), whereas the frequency chirp is performed with a voltage-controlled oscillator. The Rabi frequency profile agrees with the expected Blackman shape within a few percent. The frequency chirp, and more specifically its final frequency, was not verified as it cannot be easily checked at the required precision level of 10 Hz for HBP. The selected atoms then interact in a second microwave cavity with a 9.2 GHz microwave field synthesized from a high stability quartz oscillator weakly locked to the output of a H maser. The two  $\pi/2$  Ramsey interactions are separated by 500 ms. The number of atoms  $N_{F=3}$  and  $N_{F=4}$  are finally measured by detecting the fluorescence induced by a pair of laser beams located below the molasses region. From the transition probabilities  $N_{F=4}/(N_{F=3} + N_{F=4})$  measured on both sides of the central Ramsey fringe, an error signal is computed to lock the microwave interrogation frequency to the atomic transition using a digital servo loop.

The transition probabilities in the selection cavity are first measured as a function of the maximum Rabi frequency  $\Omega_{\max}/2\pi$ , for the Blackman and half-Blackman pulses. In this evaluation phase only, the pushing beam has to be off. To reject the fluctuations of the initial number of atoms, we measure the ratio of the atoms transferred into  $|F = 3, m_F = 0\rangle$  and the total number of detected atoms, in all magnetic sublevels. The transfer probability is then obtained by dividing the measured ratio by the initial fraction of atoms in the  $|F = 4, m_F = 0\rangle$  state. The results are shown in Fig. 2 and reproduce very well the numerical simulations.

As the maximum Rabi frequency during the experiment was set to 7.5 kHz, the resonance frequencies for transitions between  $m_F \neq 0$  states have to be significantly shifted away from the 0-0 transition. A magnetic field of  $\sim 180$  mG is applied during the pulses which keeps the total parasitic excitation of magnetic field sensitive transitions below 0.3%. This pulse induces a transient quadratic Zeeman shift on the 0-0 transition of about 14 Hz that must be taken into account to meet the resonance condition for HBP.

We alternatively measure the mean atom number for the Blackman and half-Blackman pulses during sequences of 50 fountain cycles, and compute their ratio  $R$ . We then calculate  $R_N$ , the average of  $R$  for  $N$  successive sequences. In Fig. 3, the standard deviation  $\sigma_R(N)$  for various  $N$  is plotted. The stability of  $R$  reaches  $3 \times 10^{-4}$  after a one-day integration. This reflects the insensitivity of the AP to the experimental parameter fluctuations. The mean value of the ratio is  $R = 0.506$ , whereas it was expected to be 0.5 at the  $10^{-3}$  level. This deviation cannot

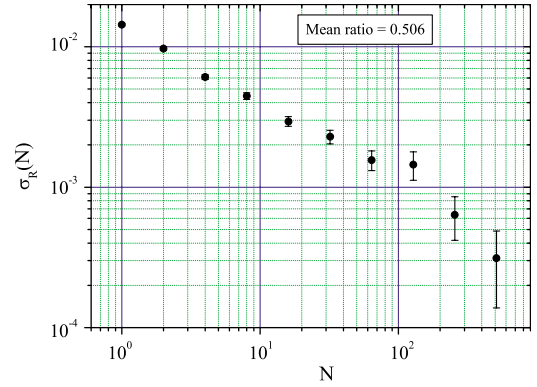


FIG. 3 (color online). Standard deviation of the fluctuation of the ratio  $R(N)$ .

be explained by a nonlinearity of the detection, which could arise from absorption in the detection beams. When the absorption in the detection laser beams is changed by a factor of 2, by changing the power of the beams, we observe no change in the ratio larger than  $10^{-3}$ . We attribute this deviation to the uncertainty in the final frequency of the sweep. In our present setup, the sweep is generated by an oscillator whose specified accuracy is limited to 50 Hz for a frequency sweep from  $-5$  to  $+5$  kHz (this difficulty can be solved by using a dedicated DDS numerical synthesizer). We measure a linear deviation in the transition probability of  $7.5(3) \times 10^{-5} \text{ Hz}^{-1}$  in agreement with the predicted value. This can explain a deviation of the ratio by about  $4 \times 10^{-3}$ . However, it is important to notice that even when the final frequency is detuned by 50 Hz, the spatial variation of the transition probability across the atomic sample is less than  $10^{-3}$ . All the tests performed here demonstrate that AP is at least accurate at the 1% level.

Measurements of the collisional frequency shift are then carried out using BP and HBP pulses. In order to amplify the collisional shift, the number of detected atoms ( $N_{\text{det}} \sim 10^7$ ) is about 10 times larger than needed in order to operate the clock in optimal conditions. From a differential measurement, one can extrapolate the frequency  $\nu_0$  of the clock at zero density with respect to the H maser. The relative resolution of the frequency difference is  $2 \times 10^{-13} \tau^{-1/2}$ , limited by the phase noise of the quartz oscillator used for the fountain interrogation. To check whether this extrapolation is correct, we measure the corrected frequency  $\nu_0$  for two different initial temperatures of the atomic cloud, 1.1 and 2.3  $\mu\text{K}$ , for which the effective densities, number of detected atoms, and  $K$  coefficients are expected to be different. We switch every 50 cycles between four different configurations: 1.1  $\mu\text{K}$  and BP, 1.1  $\mu\text{K}$  and HBP, 2.3  $\mu\text{K}$  and BP, 2.3  $\mu\text{K}$  and HBP. This rejects long term fluctuations in the experiment induced by frequency drift of the H maser used as a reference, variation of the detection responsivity, and fluctuations of other systematic effects.

TABLE I. Measurement of the cold collisional frequency shift using adiabatic passage for two atomic temperatures  $T$ . The first (second, when present) error bar indicated in parenthesis reflects the statistical (systematic) uncertainty. The difference in corrected frequency  $\delta\nu_0$  is  $-0.012(7)(5)$  mHz.

$T$	$\delta\nu_{\text{diff}}$ (mHz)	$R$	$K (\times 10^{-11})$ Hz/atom
1.1 $\mu\text{K}$	$-0.323(5)$	0.5063(3)	$-8.62(13)$
2.3 $\mu\text{K}$	$-0.260(5)$	0.5056(3)	$-10.04(20)$

The results are summarized in Table I. For each configuration, the measurement of the clock frequency and number of atoms is averaged over a total time of about 50 h. One can then extract the difference between the high and low density collisional shifts  $\delta\nu_{\text{diff}}$  with a relative resolution of  $5 \times 10^{-16}$ . The  $K$  constants are found to differ by about 20%. The corrected frequency  $\nu_0$  is then simply obtained by subtracting  $\delta\nu_{\text{diff}}$  to the measured clock frequency in the low density configuration. The difference between the corrected frequencies  $\delta\nu_0$  can then be estimated. The uncertainty on this measurement is twofold, a statistical uncertainty, and a systematic error which reflects the 1% uncertainty on the ratio  $R$ . We find  $\delta\nu_0 = -0.012(7)(5)$  mHz, which is less than 2% of the collisional shift at high density, and compatible with zero within the error bars.

To compare the standard selection method (SSM) and AP, we perform a measurement alternating both methods for selection. Table II displays the results obtained for a temperature of 1.1  $\mu\text{K}$ . The  $K$  coefficients are found to differ by about 10%. Indeed, when using a  $\pi/2$  pulse with the standard selection, the density distribution is transversally distorted: atoms along the symmetry axis of the cavity are more efficiently transferred than off-axis atoms. This increases the effective density for the same number of detected atoms with respect to AP, giving a larger collision shift at low density. In fact,  $K$  is expected to be lower with SSM than AP, in agreement with our measurement. Extrapolating the frequency to zero density when using SSM then leads to an error of about  $3 \times 10^{-15}$  at this density.

It is also important to notice that we measure simultaneously the collisional frequency shift and the shift due to cavity pulling [16], which is proportional to the number of atoms crossing the Ramsey cavity. Both are correctly evaluated by our method.

In conclusion, we demonstrate here a powerful method based on an adiabatic transfer of population to prepare atomic samples with a well-defined density ratio. An important point is that the cold collisional shift is measured precisely without any absolute calibration or numerical simulation. This holds even when parameters of

TABLE II. Comparison between the adiabatic passage technique (AP) and the standard selection method (SSM). The temperature of the sample for these measurements was 1.1  $\mu\text{K}$ .

	$\delta\nu_{\text{diff}}$ (mHz)	$R$	$K (\times 10^{-11})$ Hz/atom
SSM	$-0.234(7)$	0.540(3)	$-8.08(24)$
AP	$-0.275(8)$	0.5054(8)	$-8.97(23)$

the atomic sample, or even of the atomic detection, are fluctuating. This method can lead to a potential control of the cold collision shift at the  $10^{-3}$  level, or even better. This capability could be demonstrated by using an ultra-stable cryogenic oscillator [9], allowing a frequency resolution of  $10^{-16}$  per day. Having now at hand a powerful method to determine the collisional shift with the required resolution, the evaluation of the Cs fountain accuracy at the  $10^{-16}$  level is reachable. This technique can be readily extended to the optical domain and should be applicable to other high precision measurements using cold atoms or quantum degenerate gases for the evaluation of the phase shifts induced by atomic interactions.

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