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A simple velocity-tunable pulsed atomic source of slow metastable argon

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Abstract

A pulsed beam of metastable argon atoms having a low tunable velocity (10 to 150 m s⁻¹) is produced with a very substantial brightness (9 × 10⁸ Ar* s⁻¹ sr⁻¹). The present original experimental configuration leads to a variable velocity dispersion that can be smaller than the standard Brownian one. This behaviour, analysed using Monte Carlo simulations, exhibits momentum stretching (heating) or narrowing (cooling) entirely due to a subtle combination of Doppler and Zeeman effects.

Keywords: slow atomic beam, small velocity dispersion, metastable argon, magneto-optical trap

(Some figures may appear in colour only in the online journal)

1. Introduction

Atomic manipulation with lasers has been a laboratory standard for more than thirty years, especially through different atomic beam slowing techniques [1]. A broad range of improvements have been developed with the Zeeman slower principle to load a magneto-optical trap (MOT). Consecutively, for the MOT technique’s development, the need for colder clouds led to the idea of multi-stage trapping. The first trap is filled from a thermal source, while the second is loaded from the first source in the mK range. Then, a cold atomic beam is made from the first trap. The realisation of such cold beams required the refinement of different unbalanced MOT configurations. The first low velocity intense source (LVIS) [2] opened the way for related continuous low-velocity sources [3, 4]. Novel configuration and improvements are still in progress e.g. [5–8]. Therefore, velocity-tunable beams are rarely studied [9] and especially those ranging from 20 to 100 m s⁻¹ [10, 11].

In this paper an original method of producing a simple pulsed, low-velocity, tunable beam is presented for applications in atom interferometry experiments that require a well-defined de Broglie wavelength and critical surface purity that metastable atoms do not alter. In section 2, we describe the experimental set up, including a metastable supersonic beam, and an uncommon experimental analysis of the Zeeman’s slower performance for various operating currents. Section 3 shows both the experimental and theoretical results of the pulsed low-velocity tunable beam. Then, the ‘stretching’ (heating) or ‘narrowing’ (cooling) effect on the velocity dispersion is studied in detail in section 4. Section 5 presents a consideration about the atomic polarisation before the concluding remarks in section 6.
state argon atoms $\Phi_{\text{sat}} = 1.5(3) \times 10^{16}$ s$^{-1}$ is measured with a torsion pendulum. Argon metastability is obtained directly by electronic excitation with an electron gun (HeatWave Labs 102 245). A flux of metastable argon atoms $\Phi_{\text{Ar}} = 3 \times 10^{8}$ s$^{-1}$ is measured with a Faraday cup by secondary electron emission. The FWHM spatial divergence of 5.2 mrad is measured with a time-position detector made of two 80 mm diameter micro-channel plates (MCPs) in a chevron configuration followed by a delay-line detector (DDL80, RoentDek GmbH). The experimental brightness is consequently $1.5 \times 10^{13}$ Ar$^{*}$ sr$^{-1}$ s$^{-1}$. The relatively low excitation efficiency $\Phi_{\text{Ar}}/\Phi_{\text{sat}} = 2 \times 10^{-8}$ given by this technique, in comparison with approximately $10^{-4}$ from discharge set-up [13], is balanced by the absence of beam deflection and a transverse cooling stage before entering the Zeeman slower. Note that the basic process of metastable atom detection is the secondary electron emission from a conducting solid. This emission originates from their high internal energy, making the quantum yield almost independent of the rare gas atom under consideration. A small velocity dispersion ($\Delta v/v \sim 10\%$) and the absence of residues (ions, electrons and ultraviolet (UV) photons) allow efficient slowing conditions. Nevertheless, the high flux of argon in the ground state makes too high partial pressure for trapping. A 2 mm hole positioned at the Zeeman slower entrance is required to reduce this flux before trapping the atoms.

The experimental set up was described in detail in [14]. In brief, the Ar$^*$ cycling transition $^3P_2 \leftrightarrow ^3D_3 (\Gamma = 2\pi \times 5.58$ MHz and $\Lambda_{\text{res}} = 811.531$ nm in vacuum) is used. A $-340$ MHz red detuned laser beam is injected into the Zeeman slower made of two independent magnetic coils: the first one is fixed, whereas the second one offers the possibility of controlling the output velocity from 230 to 30 m s$^{-1}$ merely by changing its current.

Before acting the trapping set up, the time-position detector is positioned on the atomic beam axis. A mirror in between the Zeeman slower output and the detector allows the laser beam to counterpropagate with respect to the atomic beam with a power of 60 mW giving a maximum intensity of a few times $I_{\text{sat}}$, with $I_{\text{sat}} = 1.4$ mW cm$^{-2}$. A time of flight (TOF) technique is realised with a mechanical chopper at the entrance of the Zeeman slower giving both velocity and dispersion. The measurements are reported in table 1 for different velocities by changing the second Zeeman coil current. As expected, when the number of spontaneous emitted photons increases during the slowing process, the velocity dispersion increases as well towards smaller velocities, attaining its largest values (46.5%) at 33 m s$^{-1}$, showing a broad explosion of the atomic beam. Simultaneous two-dimensional (2D) beam images complete the analysis by giving the FWHM angular apertures $\Theta$, which follow the same behaviour. Additionally, the effective source position can be calculated from a Monte Carlo simulation in [14], with $R_{\text{c}} = 0.257 \lambda_{\text{ub}} D_{\text{off}}/D_{\text{s}}$ ($D_{\text{s}}$ being the distance from a pair of Young slits and $\lambda_{\text{ub}}$ being the effective source radius), leading to the transverse coherence radius, which decreases from 490 nm at 250 m s$^{-1}$ to 60 nm at 60 m s$^{-1}$.

### Table 1. Velocity dispersion over velocity ($\Delta v/v$), angular aperture ($\Theta$) and calculated coherence radius ($R_{\text{c}}$) for several velocities at the output of the Zeeman slower.

| $v$ (m s$^{-1}$) | 179 | 135 | 88 | 64 | 44 | 33 |
|----------------|-----|-----|----|----|----|----|
| $\Delta v$ (m s$^{-1}$) | 5.7 | 4.6 | 8.5 | 15.5 | 17.3 | 15.3 |
| $\Delta v/v$ (%) | 3.2 | 3.4 | 9.7 | 24.2 | 39.4 | 46.5 |
| $\Theta$ (mrad) | 13 | 19 | 35 | 44 | 41 | 51 |
| $R_{\text{c}}$ (nm) | 100 | 80 | 60 | 60 | 57 | 55 |

3. Low-velocity-tunable pulsed atomic beam

The chosen option for making the low-velocity atomic beam was to start from a standard MOT. A 6.0" spherical octagon vacuum chamber (MCF600-SphOct-F2C8, Kimbell) pumped by an Agilent’s 201 s$^{-1}$ StarrCell ion pump is placed at the output of the Zeeman slower. The lowest reachable pressure is limited to $5 \times 10^{-9}$ mbar due to the high flux of argon in the ground state arriving in the chamber. The red-detuned slowing laser is injected into the Zeeman slower vacuum tube after crossing the MOT chamber (see figure 1). Optical MOT beams are realised with three retro-reflected laser beams, 25 mm in diameter and 11$^\circ$ red-detuned, on the same closed transition as the Zeeman slower with a total power of 48 mW. Magnetic fields are generated by an anti-Helmholtz coil configuration providing an axial magnetic gradient $\partial B_z = 0.38$ G cm$^{-1}$ A$^{-1}$ and $B_z' = \partial^2 B_z = 0.2$ G cm$^{-1}$ A$^{-1}$ operating in a range of 10–60 A. A small CMOS camera allows us to estimate the number of trapped atoms ($N_{\text{MOT}} = 6 \times 10^4$) and the cloud FWHM diameter ($d_{\text{MOT}} = 100$ µm at 60 A). The loading time is 100 ms. Additional magnetic fields are used to translate the MOT in the three directions up to a few cm in order to avoid the large atomic flux from the source with no significant change. The temperature is evaluated to be in the range of a few mK. The slow atomic beam is thus produced perpendicular to the source beam in a large UHV chamber terminated by the time-position detector at 1.1 m from the trapped cloud.

The production of the slow atomic beam follows the Zeeman slower principle in an acceleration configuration. The applied force comes from the recoil velocity, $v_{\text{rec}} = \hbar k/m = 12.3$ mm s$^{-1}$ for a single absorbed photon. For a complete phenomenon description the scattering rate $\Gamma_{\text{sc}}$ of the atoms in a field of intensity $I$ will be decomposed into $\Gamma_{\text{abs}}$, the absorption rate, and $\Gamma$, the emission rate, (e.g. in [15]) with:

$$\Gamma_{\text{sc}} = \frac{\Gamma}{2} \left( \frac{s_0}{1 + s_0 + (2\Delta/I)^2} \right)$$

and

$$\Delta = (2\pi)\delta I - k_L v + (g_f m_f - g_i m_i)\mu_B B/h,$$

where $\delta I = c(k_L - k_{\text{res}})/(2\pi)$ is the detuning in MHz, $k_L = 2\pi/\lambda_L$ is the laser wavenumber, $k_{\text{res}} = 2\pi/\lambda_{\text{res}}$ is the Ar$^*$ resonance wave number, $c$ is the speed of light, $s_0 = I/I_{\text{sat}}$ is the saturation parameter, $\mu_B$ is the Bohr magneton, $g_f = 1.338$,
$g_1 = 1.506$ is the Landé factors, $m_f = 3$, $m_i = 2$ and $B$ are the magnetic field strength. The atomic motion and pushing laser are assumed to be co-propagating in (2), \( \vec{k}_i \cdot \vec{v} = k_i \nu \), which is obviously the case in the described experiment. According to \( \langle t_{ac} \rangle = \langle t_{abs} \rangle + \langle t_{em} \rangle \) (see [14] for a complete demonstration) \( \Gamma_{abs} \) can be defined as:

\[
\frac{1}{\Gamma_{abs}} = \frac{1}{\Gamma_{sc}} - \frac{1}{\Gamma} \tag{3}
\]

Starting from trapped atoms in an MOT with velocities in the range of cm s\(^{-1}\), the mean atomic beam velocity after an irradiation time \( t_p \), the pushing time, is determined by the number of absorbed photons. The entire mechanism for adjustable velocities takes necessarily into account all the possible parameters involved: pushing laser intensity, detuning, time \( I, \delta_1 \) and \( t_p \), the Doppler effect, optical molasses forces and magnetic field strength \( B \), if any. As the objective is to cover a broad range of small velocities with a well-known velocity dispersion, the choice of pulsed atomic regime has been made for accurate TOF measurements. Such a time pulse is calculated to be smaller than 3 ms, to get \( v_{\text{max}} = 150 \text{ m s}^{-1} \), for a laser intensity close to saturation and detuning in a range of a few \( \Gamma \). In this small range of time, the MOT coils are left on to avoid transient uncontrolled magnetic fields but the optical MOT beams are turned off with an AOM. Keeping the magnetic field on allows us to reach a relatively high repetition rate, 12 Hz, for the MOT cloud to be recovered. A second AOM drives the time sequence \( t_p \) and the whole routine is Labview\textsuperscript{TM} controlled. The pushing laser power is 7 mW with a 4.4 mm waist giving \( I = 12 \text{ mW cm}^{-2} \) and \( s_0 = 8 \). The goal here is different from that of atomic fountains: it is to generate a beam with a tunable and well-defined velocity within a limited angular spread. As explained below (see section 4) this goal has been reached. The slow atom beam is characterised by arrival time-positions and the number of atoms recorded after 1 m free flight as the distance range of acceleration is less than 10 cm. It is first noted that the atomic cloud cannot completely pass the magnetic barrier for pushing times shorter than 0.3 ms with \( B_z' = 12 \text{ G cm}^{-1} \) (60 A) (figure 2). The brightness of the atomic beam is also evaluated with atomic positions on the detector. A quantum efficiency of 15% for the MCPs and \( \text{Ar}^* \) is assumed (see, e.g. [16]). An increase in the brightness with velocity is shown. The saturation in figure 2 corresponds to a constant velocity while increasing \( t_p \).

To get a final velocity high enough to pass the magnetic barrier (20 m s\(^{-1}\) for \( B_z' = 12 \text{ G cm}^{-1} \)), the remaining parameters that strongly influence the final velocities for a given \( t_p \) are the pushing laser detuning \( \delta_1 \), and the MOT magnetic field strength \( B \). Figure 3 presents the final velocity over \( t_p \) for various detuning and magnetic trap intensities. A broad range of velocities are attainable (from 10 to 150 m s\(^{-1}\)) by adjusting respectively the above parameters. In such a short range of pushing time, it is crucial to keep \( \Delta \) as small as possible during the acceleration due to the competition between the Doppler effect, the Zeeman effect and laser detuning. The lines (red) are calculated with true experimental parameters (magnetic field intensity, pushing laser detuning and intensity) and take into account the local \( \Gamma_{abs} (B, \nu) \). \( B(z) \) follows the measured magnetic gradient \( (B(z) = 0.2 \text{ G cm}^{-1} \text{ A}^{-1} \) until 6.2 cm from the cloud position, before it decreases with the most appropriate theoretical gradient \( B(z) = -0.04 \text{ G cm}^{-1} \text{ A}^{-1} \).

**Figure 1.** Experiment representation. Atoms (\( \text{Ar}^* \)) come from the right-hand side through the second part of the Zeeman more slowly; the MOT chamber occupies the central part; a large experimental chamber is placed on the left-hand side; the slow atomic beam follows the pulsed pushing laser beam direction represented by the blue arrow (\( z \)-axis). The time-position detector stands on the far left of the experimental chamber.

**Figure 2.** \( \text{Ar}^* \) number \( N \) (blue circles) and brightness \( Br \) (red squares) for various pushing times. \( \delta_1 = 2.1 \Gamma, B_z' = 12 \text{ G cm}^{-1} \) and repetition rate is 4 Hz.
Figure 3. Velocity as a function of pushing time $t_p$; the points are experimental data and the lines come from a simulation with no adjustable parameters. Three different pairs of parameters are shown: (red open circles: $\delta_L = 26$ MHz; $B'_x = 2$ G cm$^{-1}$), (black squares: $\delta_L = -4$ MHz; $B'_x = 12$ G cm$^{-1}$) and (green dots: $\delta_L = 28$ MHz; $B'_x = 12$ G cm$^{-1}$).

used in the model. The excellent agreement, found with no adjustable parameters, confirms the model’s robustness for the average velocity over the pushing time. For a chosen triad of the major experimental parameters ($t_p$, $B$, $\delta_L$), the mean velocity is extremely stable and reproducible over weeks; laser intensity could be introduced as a parameter but with no substantial gain and being careful to keep $s_0 > 1$. In a desired range of velocities, $B$ and $\delta_L$ might be anticipated, and the pushing time $t_p$ parameter allows easy and fast velocity changes. The maximum acceleration distance is found to be less than 20 cm for large velocities, making this technique compact enough for most interferometry experiments.

4. Stretching or narrowing of the velocity dispersion

The initial atomic velocity in the MOT is very low compared to that induced by the action of the pushing laser, i.e. by the absorption and re-emission of a large discrete number of photons. This is the reason why the velocity distribution is not of the Boltzmann type but rather Poissonian. The present experimental configuration demonstrates a subtle effect that exhibits stretched or narrowed velocity dispersions with respect to the free Brownian motion. In the following, the laser detuning and magnetic field are set with three different pair values: ($\delta_L = 26$ MHz; $B'_x = 2$ G cm$^{-1}$), ($\delta_L = -4$ MHz; $B'_x = 12$ G cm$^{-1}$) and ($\delta_L = 28$ MHz; $B'_x = 12$ G cm$^{-1}$). The averaged velocity $\bar{v}$ over the pushing time $t_p$ is well described (see above) and the velocity dispersion analysis requires the separation of the transverse ($x$, $y$) and longitudinal ($z$) directions as independent phenomena. Two independent models will be used to calculate the three components of the velocity dispersion. The first method calculates the expected Brownian dispersion in its configuration by considering only the final number of absorbed photons during the entire acceleration. The second method is a complete Monte Carlo calculation, which was fully developed in [14], for a Zeeman slower analysis and is easily adapted to the present configuration. It decomposes each absorption–emission process during the pushing sequence as a unique random trajectory with the absorption time $t_{ab} = -\log[Rn(0,1)]/T_{abs}$ and the emission time $t_{em} = -\log[Rn(0,1)]/T$ with $Rn(0,1)$ a random number in [0,1]. The atomic velocity is changed by an additional momentum $\hbar k$ along $z$ for one absorbed photon and the photono- emission follows the emission diagram in $(\cos^2 \theta + 0.5 \sin^2 \theta)$ for the laser light polarised $\sigma^+$ with $\theta$ the spherical angle over the $z$ axis (see e.g. [15]). This sequence leads to a number of absorbed–emitted photons $N_i$ and a unique final velocity $\bar{v}_i$. A series of a few hundred independent trajectories is usually enough to build up a significant distribution with the average number of absorbed photons $\bar{N}$ and $\bar{v} = \bar{N}v_{rec}$. The Q-Mandel’s parameter $Q = -2\Omega^2(3\beta^2 - \delta^2_1)/(2\delta^2_L + 2\beta^2 + \Omega^2)$ [17, 18] is in between $-0.6$ at resonance down to 0.01 for $\Delta_L = \pm 50$ MHz with $\beta = 2.62$ MHz (half the Einstein coefficient $A$ for the $\Omega_2^2D_2$ transition in argon) and the Rabi frequency $\Omega$ about 5 MHz. The scattering-like distribution appears to be a pure Poisson distribution almost everywhere, whereas it is quasi-sub-Poissonian at resonance. In the presented configuration atoms can effectively be at resonance for a very short time with regard to the chosen parameters. This will play a major role in the velocity distribution.

As demonstrated below, this approach is a unique way to reproduce precisely the longitudinal velocity dispersion measured in such a complex configuration.

4.1. Transverse velocity dispersion

Figure 4 shows the angular dispersions FWHM recorded on the detector image, vertically ($y$) and horizontally ($x$). As the atoms are free flying after the pushing sequence, the angular dispersion is directly linked to the velocity distribution in the transverse plane. The dashed black line is the analytical Brownian dispersion calculation assuming independent iterative processes of $\bar{N}$ events leading to a standard deviation in $\sqrt{\bar{N}}$. Nevertheless, all random emissions are in three-dimensional (3D) space and a $\sqrt{3}$ factor is requested.
for a correct standard deviation calculation. An additional corrective factor ($\alpha_{v,z} = 0.95$) to the spherical emission is added to include the photon emission diagram integration over the perpendicular plane. Then the standard deviation becomes $\sigma(\bar{N}) = \alpha_{v,z} \sqrt{\bar{N}/3}$ (see e.g. [19]). With $v = \bar{N}_{\text{rec}}$, the angular dispersion is finally written as:

$$\sigma_\theta = \arctan\left(\frac{\sigma(v)}{\nu}\right) \approx \frac{\nu_{\text{rec}}}{\nu}$$

(4)

A complete Monte Carlo calculation (not shown in the figure) perfectly fits the Brownian curve. A good agreement is found with the experimental data $\sigma_\theta$, at high velocities. A classical calculation shows that the magnetic gradient has an impact on horizontal trajectories, especially at low velocities or large magnetic gradient, which explains the slight difference between $\sigma_\theta$ and $\sigma_\theta$. Finally, it is demonstrated that the transverse dispersion is only determined by the final number of absorbed photons $\bar{N}$, as a pure Brownian motion.

4.2. Longitudinal velocity dispersion

Contrary to the transverse velocity dispersion, the longitudinal velocity dispersion $\sigma_v$ is governed by two independent, but mixed, random phenomena: the dispersion of the absorbed photons number $N_{z_{\nu}}$ through $v = \bar{N}_{\text{rec}}$, after the finite pushing time $t_\nu$ and the spontaneous emission dispersion along $z$.

The measurements were made under two experimental conditions: (i) with a large initial detuning ($\delta_{L} = 26 \text{ MHz}$) and a low magnetic gradient ($B'_z = 2 \text{ G cm}^{-1}$) aimed at producing low velocities (black squares in the upper part of figure 5), (ii) with an initial detuning close to resonance ($\delta_{L} = -4 \text{ MHz}$) and a stronger magnetic gradient ($B'_z = 12 \text{ G cm}^{-1}$), which favours the production of high velocities (red full circles in the same figure). It can be seen that $\sigma_v$ considerably varies, in a way strongly dependent on the experimental conditions.

Indeed, in this situation the interpretation is much more complex than before because of the spatial dependence of $B$ and of the Doppler effect, which consequently change $\Gamma_{\text{abs}}(z,v)$ along the acceleration process. The key question is to determine the number of absorbed photons, equal to that of the emitted photons, during $t_\nu$ for each series of random times used for absorption (which are $z$ and $v$ dependent) and for emission.

In a first attempt a purely Brownian approach is used. In this approach only the last local value of $\Gamma_{\text{abs}}(z,v)$ is considered rather than its full history as the Monte Carlo calculation does. This Brownian dispersion takes into account the local $\Gamma_{\text{abs}}(z,v)$ and $\alpha_z = 1.11$ the corrective coefficient of the photon emission diagram to spherically symmetric emission. As demonstrated in appendix (A.4) and with the standard deviation composition rule, the standard deviation of the longitudinal velocity $\sigma_v$ can be written as:

$$\sigma_v = \nu_{\text{rec}} \sqrt{\frac{\Gamma_{\text{abs}}^2 + \Gamma^2}{(\Gamma_{\text{abs}} + \Gamma)^2} + \frac{\alpha_z^2}{3}}$$

(5)

In figure 5, equation (5) is plotted (FWHM) in solid dark and dashed red for the two different experimental series (i) and (ii) defined above. It can be noted that the complete process calculation is always smaller when far off resonance ($\Gamma_{\text{abs}} \rightarrow 0$ and $\sigma_v \rightarrow \nu_{\text{rec}} \sqrt{\bar{N}/3}$) than a full process kept at resonance ($\sigma_v = \nu_{\text{rec}} \sqrt{\bar{N}/0.5 + \alpha_z^2/3}$). In situation (i), compared to the Brownian model, $\sigma_v$ is first larger (momentum stretching) at velocities below 45 m s$^{-1}$, then smaller (momentum narrowing). In situation (ii), $\sigma_v$ is first smaller, then larger and finally smaller again beyond 115 m s$^{-1}$.

The Monte Carlo approach finds all its necessity in such a specific configuration, to model what the Brownian approach failed to explain. Indeed, the entire history of all possible atomic paths has to be taken in account. For the two series, a qualitative agreement of the dispersion evolution is plotted with up-triangles and down-triangles. In both situations, the Monte Carlo results reproduce the shape, although with an off-set below the experimental data that demonstrates the existence of an extra heating phenomenon. Figure 5 (lower part) shows for the same couples of parameters the detuning, which is a function of $v$ and $z$, over the final velocity. There are obvious correlations between high-velocity dispersion and very small detuning. Conversely, there is strong decrease in the velocity dispersion where the detuning is large.
The momentum compression can be attributed here to different atom histories. The slowest atoms get smaller $I_{\text{abs}}(z, v)$ than faster atoms. Then, slower atoms absorb a larger number of photons than the faster atoms do, leading to a momentum compression.

The narrowing phenomenon presented above can be compared to what happens during the Zeeman slowing process. The Zeeman slower is built such that the atoms are as close as possible to resonance in order to minimise the slowing length. The longitudinal velocity distribution has been poorly studied [20, 21], but it appears to be larger than that predicted by a simple Brownian model. In the present experimental configuration, the larger part of the accelerated process can be chosen close to resonance before having a momentum focalisation far off resonance.

An additional filtering can be implemented to reduce the velocity dispersion. The atomic path, in the described experiment, is 1.1 m leaving a free area for a synchronised resonant laser beam perpendicularly aligned. By a fine tuning of the shutters sequence, it is possible to blow away the faster atoms, those positioned on the head of the atomic cloud, at a high laser intensity (a few tens of $I_{\text{sat}}$). The superposed TOFs are shown in figure 6 with and without an optical chopper at an arbitrary chosen velocity of 45 m s$^{-1}$ with ($\delta_{L} = 14$ MHz, $B'_{z} = 6$ G cm$^{-1}$). Note that a continuous atomic beam has been made with a brightness up to $2.5 \times 10^{9}$ s$^{-1}$ . sr$^{-1}$, the velocity of which is accessible by TOF, e.g. using an accessory chopping laser beam if requested.

5. Atomic polarisation

In many experiments, atoms with a magnetic dipole moment, such as metastable $^1P_2$ argon atoms, are used. Because of the dependence on the magnetic number $M$ of various effects (e.g. the effect of external magnetic or radiation fields, or effects due to the vicinity of a solid surface), it is important to consider the spin evolution in between the source to some point at a distance typically of a few tens of centimeters where a scattering object is placed. Starting from a fully polarised ($M = +2$ referred to the $z$-axis) atom beam generated by the $\sigma^+$ polarised pushing laser beam, the spin dynamics in the presence of an ad hoc guiding magnetic field has been examined, using the time-dependent Schrödinger equation. In this zone, the magnetic earth field is reduced by a magnetic shielding; a longitudinal field (10 $\mu$T along the $z$-axis) is added to the fringe field of the MOT and an additional transverse constant field parallel to $y$ (maximum amplitude of 0.5 mT) is applied in front of the target. Under these conditions, it is found that for velocities in the range 35–45 m s$^{-1}$, almost only the $M = +2$ final state (referred to the $y$-axis) is significantly populated.

6. Conclusion

We obtain a metastable argon beam with a tunable velocity ranging from 10 m s$^{-1}$ up to 150 m s$^{-1}$. Its characteristics in terms of flux, velocity dispersion and angular aperture make it readily usable in a wide variety of scattering experiments in a large sense, involving external fields, structured surfaces, graphene foil, well-localised electromagnetic fields etc. An experimental and theoretical demonstration of momentum stretching or narrowing has been achieved. These original properties originate purely from the magnetic field topography with regard to Doppler effect. A programmable power supply for the pushing laser AOM to compensate the magnetic field (chirp-cooling [1]) would lead to much better control of the atomic cloud properties, and especially a substantial reduction in the pushing time and consequently a net decrease in the full cycle time. Owing to general considerations about the complete acceleration process, the momentum statistics are discussed in detail with regard to the not so-studied Zeeman slower dispersion mechanism. An investigation of the spin time dynamics is realised in the presented experimental conditions. Despite the need for trapping, such a simple atomic source gets better properties than a simple Zeeman slower.

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Appendix

Let us start from the fact that the time needed for $N$ absorption/emission cycles follows (for large values of $N$, here more than a few hundred) a normal distribution with a mean value $N(\tau + \tau')$ and a variance $N(\tau^2 + \tau'^2)$ with $\tau = 1/I'_{\text{exp}}$ and $\tau' = 1/I'_{\text{abs}}$.

$$P(t|N) = \frac{1}{\sqrt{2\pi N(\tau^2 + \tau'^2)}} \exp \left[ -\frac{(t - N(\tau + \tau'))^2}{2N(\tau^2 + \tau'^2)} \right]$$

(A.1)
The number of cycles within a duration $t$ is in fact derived from Bayes theorem:

$$P(N|t) = \frac{P(N)}{P(t)} \frac{P(t|N)}{P(t)} = \frac{P(N)}{P(t)} \exp\left[-\frac{(N - \bar{N})^2}{2N N^2 + \tau^2}\right]$$

where $P(t|N)$ and $P(t)$ are the probability densities, and where we have introduced the mean photon number $\bar{N} = t(\tau + \tau')$. Let us assume a normal distribution for $P(N|t)$ through the approximation $N \approx \bar{N}$:

$$P(N|t) \approx \frac{P(\bar{N})}{P(t) \sqrt{2\pi N (\tau^2 + \tau'^2)}} \exp\left[-\frac{(N - \bar{N})^2}{2N N^2 + \tau^2\tau'^2}\right]$$

(A.3)

This leads directly to the desired result concerning the variance $\sigma_2$ of the fluctuations of the number of absorption/emission cycles:

$$\sigma_2 = \bar{N} \frac{\tau^2 + \tau'^2}{(\tau + \tau')^2} = \bar{N} \frac{\Gamma^2 + \Gamma_{abs}^2}{(\Gamma + \Gamma_{abs})^2}$$

(A.4)

Let us note that the normalisation of (A.3) implies $P(N) \approx (\tau + \tau')^2 P(t)$, which is a quite intuitive result as the number $\Delta N$ of cycles during $\Delta t$ is approximately $\Delta N = \Delta t/(\tau + \tau')$, so that $P(N|t) \Delta N \approx P(t) \Delta t$. To conclude, a convolution between the variance of the number of absorption/emission cycles ($\sigma_2$) and the emission ($\sigma_{em}$) processes is calculated as $\sigma_v^2 = \sigma_0^2 + \sigma_{em}^2$ with $\sigma_{em} = \nu_{rec} \sqrt{\bar{N} \alpha_j^2 / 3}$, and $\sigma_0 = \nu_{rec} \sigma$.

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