Time-delayed intensity–interferometry of the emission from ultracold atoms in a steady-state magneto-optical trap

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Abstract

Time-delayed intensity–interferometry (TDII) measurements of the fluorescent emission from an ultracold ensemble of thermal $^{87}$Rb atoms in a steady-state magneto-optical trap are presented, which reveal the underlying coherent and incoherent dynamics of the atoms. Measurements carried out with a 5 ns time resolution yielded a second-order intensity correlation function with the theoretically predicted value of 2 at zero delay. In addition coherent Rabi oscillations were seen for up to five full periods—much longer than the spontaneous emission lifetime of the excited state of Rb. The oscillations were damped out by $\sim$150 ns, and thereafter an exponential decay observed, from which the mean velocity of atoms and thus, the temperature of the ensemble was estimated. The values so obtained compare well with those determined by standard techniques. It is seen that TDII permits a quantitative study of the coherent and incoherent processes, even in a large ensemble of independent atomic emitters in random thermal motion. This measurement of second-order correlation powerful technique can reveal hidden periodicities such as coherent Rabi oscillations that are not directly seen in the emission from a large collection of atoms. In addition it can also reveal information about the mean velocity of the thermal ensemble of emitters, even though the Doppler broadening of emission due to the motion of atoms is smaller than the natural linewidth and is not directly measurable.

Keywords: time-delayed intensity–interferometry, magneto-optical trap, photon bunching

(Some figures may appear in colour only in the online journal)
D1 and D2 at locations $r$ and $r + R$ at times $t$ and $t + \tau$ and the angular brackets denote time averaging. It is the correlation of the number of photons, or intensities, that is examined, as opposed to a correlation of amplitudes in conventional interferometers; hence the name time-delayed intensity–interferometry (TDII). It has been theoretically shown (see, for example, [2]) that the second-order correlation function $g^{(2)}(\tau = 0) = 2$ for a thermal state, implying a tendency for the bunched or correlated emission of photons. It has unit value for a coherent source, indicating emission of photons at random times. For an n-photon Fock state, it equals $1 - \frac{1}{n}$ signifying anti-bunching. The value of $g^{(2)}(\tau)$ for all sources, however, approaches unity for long time delays.

Temporal bunching of photons from thermal sources, ever since the postulation of the concept, has been an intriguing phenomenon, and has been the focus of numerous experiments. The earliest laboratory thermal source studied was a Hg vapour lamp [7], light from which showed a meagre bunching of 1.17. Martiensen and Spiller devised a method of creating pseudo-thermal light by transmitting coherent laser light through a rotating ground glass plate [8] such that the time-varying surface inhomogeneities introduced temporal and spatial decoherence. In recent years, laboratory control and measurement techniques have enabled the creation of pseudo-thermal light sources with theoretically expected values of $g^{(2)}(\tau) = 2$ [9–12].

We report here TDII measurements on light from another source of bunched photons—an ensemble of laser cooled atoms. Though laser cooled atoms have been available for more than three decades, there have been very few direct TDII measurements of their emission [13–15], most of which, due to insufficient time resolution, observed only limited bunching. All measurements hitherto have been carried out in optical molasses, which were either periodically or continuously loaded with atoms precooled in a magneto-optical trap (MOT). In this paper we present measurements of the second-order correlation function of light emitted by ultracold atoms in a steady-state magneto-optical trap, where the cooling and repumper beams, and also the quadrupolar magnetic field, are kept on. We observe the ideal value of 2 for the zero-delay intensity–intensity correlation function. Damped Rabi oscillations are observed for time delays up to $\sim 150$ ns, and an exponential decay for longer time delays. Despite the fact that the emission being studied is from a collection of uncorrelated atoms that are in random thermal motion, and that the observations are averaged over an 8 h period, coherent effects are seen, bringing out the power of higher-order correlations in revealing hidden periodicities and providing a measure of coherent and incoherent dynamics. The exponential decay at long time delays was used to determine the temperature of the ensemble.

TDII was performed on the fluorescent emission from $^{87}$Rb atoms cooled and trapped in an MOT (figure 1), which differed from usual MOTs, in that the two pairs of beams in the x–y plane were steeply inclined to each other, enclosing an angle of $55^\circ$ rather than the usual $90^\circ$ so as to accomodate, within the chamber, a pair of lenses of short-focal length and high numerical aperture. These lenses were positioned facing each other, such that their focal points coincided with the center of the MOT and could thus be used to focus light onto the MOT, or collect light emitted by a small volume within the cold cloud. To avoid clipping at the lens mounts leading to undesired scattered light, the diameters of the beams in the x–y plane were restricted to $\sim 1.5$ mm while the z-beam had a diameter of 8 mm. The cooled and trapped atoms were viewed using a charge-coupled device (CCD) camera, and their number estimated by collecting part of the fluorescent light onto a femtowatt laser. The typical cloud was roughly ellipsoidal with a mean diameter of $\sim 400$ mm, and contained about 20000 atoms.

An intensity–interferometer was formed using a fiber splitter, where the two output ends were connected to two high speed avalanche photo diodes (APDs) that had quantum efficiencies of $\sim 90\%$ and a dead time of 30 ns. The output of the APDs was fed to a homebuilt, field programmable gate array-based time-tagged-single-photon recorder, Avalanche Photodiode Optical Data Acquisition System (APODAS) [16], which utilized high speed ethernet connectivity and stored, in a PC in realtime, the arrival times of all detected photons with a temporal resolution of 5 ns. Post-processing by software enabled the determination of $g^{(2)}(\tau)$ for all $\tau$ from a single recording of the data [12]. As we worked in the photon-counting mode, the expression for $g^{(2)}(\tau)$ in terms of coincidences is [14]

$$g^{(2)}(\tau) = N_1T/(N_1N_2\tau_c)$$ (1)

where $\tau$ is the delay between arrival at the two detectors, $N_1$, $N_2$ and $N_c$ are the number of counts at detector 1, detector 2 and the coincident counts respectively. $T$ is the total observation time and $\tau_c$ the time window for coincidence (arrival of two photons is considered simultaneous if they are detected within a time gap of $\tau_c$).

For determining the bunching characteristics of emission from an ultracold atomic ensemble, $^{87}$Rb atoms were laser
cooled from close to room temperature, to \(\sim 100 \ \mu\text{K}\). The MOT was extremely stable, with the lasers locked and the cold cloud in a steady-state for days. A typical run of the experiment lasted 8 h. The cold cloud was obtained, and the cooling and repumper beams, and the magnetic field were kept switched on for the entire duration of the experiment. The cold ensemble was constantly monitored by imaging it on a camera, and also by measuring the fluorescence on a femtowatt detector (see figure 1). For the purpose of determining its second-order correlation function, fluorescent light from the central region of the cloud was collected by the high-numerical aperture lens placed inside the vacuum chamber, a few millimeters from the trap center. Care was taken to ensure that no part of the laser light entered this lens, either directly, or upon being scattered by the parts of the MOT chamber. The light was conveyed by a series of mirrors and lenses to the input of the TDII setup. The count registered in the presence of the cold cloud was in the range 40000–80000 s\(^{-1}\) while in the absence of the cold cloud it reduced to \(\sim 1200\) s\(^{-1}\), confirming that it is predominantly light from the cold atoms that enters the TDII setup. Photon arrival time data were recorded for various detunings of the cooling beam. For each detuning the number of atoms trapped was estimated from the fluorescent intensity recorded on the femtowatt detector. While the total number of atoms collected ranged from \(\sim 8000\) to \(\sim 22000\), the high numerical aperture lens employed for accepting light for TDII measurements restricted the collection of light to that from approximately one-tenth of the volume of the cold ensemble. The temperature of the collection of atoms was determined by the trap oscillation method [17], and was found to range from 200 \(\mu\text{K}\) to 50 \(\mu\text{K}\) for detunings of the cooling laser varying from \(-12\) MHz to \(-22\) MHz. The effective Rabi frequencies ranged from 25 MHz to 40 MHz.

There has been skepticism about being able to obtain a good measure of bunching, and about being able to see coherent effects in TDII measurements from a collection of large numbers of atoms (see, for example, [18]). In experiments so far, several factors have contributed to the difficulty in observing the theoretically predicted behavior. All researchers have stressed the need for good timing resolution; this, in our experiment is 5 ns and is found to be sufficient to observe ideal bunching. The small time bins necessitate longer acquisition times to obtain good statistics, making the experiment long and tedious. Light has to be collected over a single spatial coherence region, contributing to further reduction in photon counts. These factors require the experimental setup to be extremely stable, and the conditions repeatable over nearly 10 h. Another factor known to degrade the observation of bunching is the presence of a magnetic field [14], because of which all measurements hitherto had atoms cooled in an MOT and then transferred to a molasses, either by switching off the magnetic field, or by transporting the atoms to another vacuum chamber. Our experiment, however, avoids these complications, and has been carried out directly in an MOT, with all cooling and repumper beams and the quadrupolar magnetic field present. Good mechanical isolation of the setup and temperature stability of the environment ensured that the lasers remained locked for the entire duration of the experiment. Constant monitoring of the MOT fluorescence allowed for corrective measures, which, however, were not required. The diffraction-limited collection lens placed within the MOT chamber, in close proximity to the cold atoms, and the subsequent spatial filtering enabled us to collect light from a small region of the MOT, over which the magnetic field was uniform within 2 mG, eliminating broadening due to Zeeman shift. Likewise, the low temperature and thus the sub-natural Doppler width ensured that the Rabi frequency was the same for all atoms. The small size of the cold cloud (400 \(\mu\text{m}\) across) and the low number of atoms ensured that reabsorption of the emitted light was negligible. This allowed us to detect coherent effects such as Rabi oscillations.

Two-time intensity–correlation values were derived from the TDII measurements. Figure 2(a) shows the second-order correlation functions of the input laser light determined before it enters the MOT chamber, and of the light emitted by the ultracold atoms trapped in the MOT. It is clear that absorption and re-emission by the thermal ensemble of atoms has led to bunching in light, which was initially coherent \((\varphi_{21}(\tau) = 1)\) for all \(\tau\). Several interesting features are observed in \(g^{(2)}(\tau)\) for light from the cold atoms. Periodic oscillations, are seen; these are damped out by \(\sim 150\) ns, and thereafter the curve decays steadily towards unit value. The curve \(g^{(2)}(\tau)\) is less noisy for small \(\tau\) than for larger ones. As the cooling beams are further detuned from the \(5S_{1/2}F = 2 \rightarrow 5P_{3/2}F = 3\) transition (the so-called ‘cooling transition’), the value of \(g^{(2)}(0)\) approaches closer to 2 while the oscillations in correlation are more prominent and more rapid, and their damping is slower (figure 2(b) and (c)).

These periodic variations are the coherent Rabi oscillations that manifest themselves in second-order correlation measurements, as may be understood in simple physical terms. Assuming the Rb atom to be a two level system, irradiation by coherent, near resonant light causes two processes to occur. One is the periodic absorption and coherent emission of radiation leading to an oscillation between the excited and the ground state at the Rabi frequency—a rate determined by the intensity and detuning of the incoming radiation. The other is the absorption and spontaneous random emission, at a rate that falls exponentially with a characteristic lifetime, which, for the excited state of \(^{87}\text{Rb}\) is \(\sim 28\) ns. The atoms in the ensemble, being uncorrelated, do not emit in unison, and thus no periodicity will be evident in the direct observation of emission from the collection. However, all atoms undergo Rabi oscillations at the same frequency, and thus have a high probability of emission at the same regular interval. The second-order correlation, which is a measure of the probability of emission at time \(t + \tau\) conditioned on an emission having occurred at time \(t\), will therefore exhibit periodic maxima at regular intervals \(\tau_{R}\), the inverse of the Rabi frequency. Thus, the two-time intensity correlation measurement is a simple yet powerful technique that can reveal hidden
periodicities. The periodic oscillations arising from coherent dynamics under steady-state driving fields, nevertheless, show decay as decoherence sets in due to spontaneous emission and inter-atomic collisions. This reasoning also explains why $g^{(2)}(\tau)$ is less noisy for short time delays than for larger delays. The Rabi frequency in our experiment being a few tens of MHz, coherent oscillations should occur at the time scales of a few tens of nanoseconds. The lifetime of the excited state is $\sim 28$ ns, implying that a decay in amplitude of oscillation will occur over $\sim 100$ ns (a few lifetimes). Interatomic collisions occur at yet longer time scales. Thus, short delays show cleaner curves.

Considering a thermal collection of $N$ independent atoms under the action of a near-resonant driving field, the electric field $\mathcal{E}(\vec{r})$ at a point of observation $\vec{r}$ arising as the resultant of coherent emissions from individual atoms, $j$, is given as $E(\vec{r}) = \sum_{j=1}^{N} \mathcal{E}_j \sim \sum_{j=1}^{N} K_j S_j(e \rightarrow g)e^{i\phi}$, where $S_j(e \rightarrow g)$ is the de-excitation operator for the two-level atom, and $\phi_j$ the phase of the electric vector [19]. Using this in the definition of the second-order intensity correlation function, $g^{(2)}(\tau) = \langle \mathcal{E}^\dagger(0)\mathcal{E}(\tau)\mathcal{E}(0)\mathcal{E}(\tau)\rangle / \langle \mathcal{E}^\dagger(0)\mathcal{E}(0)\rangle^2$ yields terms of the forms (a)–(f) given below:

(a) $\sum_{j=1}^{N} \langle \mathcal{E}_j^\dagger(0)\mathcal{E}_j(\tau)\mathcal{E}_j(\tau)\mathcal{E}_j(0)\rangle$

These represent single-atom contributions. At $\tau = 0$ these show antibunching—an atom cannot emit more than one photon at a time.

(b) $\sum_{i=1}^{N} \sum_{j=1}^{N} \langle \mathcal{E}_j^\dagger(0)\mathcal{E}_j(\tau)\mathcal{E}_i(\tau)\mathcal{E}_i(0)\rangle$

As the emitters are a thermal collection of uncorrelated atoms, the operators in these terms may be re-ordered and factorized to yield:

(c) $\sum_{i=1}^{N} \sum_{j=1}^{N} \langle \mathcal{E}_j^\dagger(0)\mathcal{E}_j(\tau)\mathcal{E}(\tau)\mathcal{E}(0)\rangle$

$= N(N - 1) \frac{I^2}{N^2}$, where $I$ is the total intensity at the detector due to $N$ atoms.
These are the auto-correlation terms.

\[
\sum_{i=1}^{N} \sum_{j=1, j \neq i}^{N} \left\{ E_i(0) E_j(\tau) E_j(\tau) E_i(0) \right\}
\]

These terms are related to the anomalous correlation, which, for a thermal cloud, vanish on time averaging:

\[
\sum_{i=1}^{N} \sum_{j \neq i}^{N} \sum_{k=1, k \neq i, j}^{N} \sum_{l=1, l \neq i, j, k}^{N} \left\{ E_i(0) E_j(\tau) E_k(\tau) E_l(0) \right\}
\]

and various permutations. Due to the random phases, these terms drop out on time averaging.

\[
\sum_{i=1}^{N} \sum_{j=1, j \neq i}^{N} \sum_{k=1, k \neq i, j}^{N} \sum_{l=1, l \neq i, j, k}^{N} \left\{ E_i(0) E_j(\tau) E_k(\tau) E_l(0) \right\}
\]

These too do not survive on time averaging, because of the random phases.

From these, we obtain

\[
\left\{ E_1(0) E_i(\tau) E_j(\tau) E_i(0) \right\} = N \left[ \frac{\left\{ E_1(0) E_i(\tau) E_j(\tau) E_i(0) \right\}}{I} \right]^2
\]

and finally,

\[
g^{(2)}(\tau) = \frac{\left\{ E_i(0) E_j(\tau) E_j(\tau) E_i(0) \right\}}{N T^2}
\]

\[
+ \left( 1 - \frac{1}{N} \right) \left[ 1 + \left\{ E_i(\tau) E_i(0) \right\} \right] \left\{ \frac{\left\{ E_i(0) E_i(\tau) \right\}}{I} \right\}^2.
\]

where \( I = \frac{1}{N} \) is the intensity due a single atom. In the limit of a large number of atoms, the first term in equation (2), which represents the single-atom contributions, becomes insignificant, giving way to the standard relation between the first- and second-order correlations: \( g^{(2)}(\tau) = 1 + |g^{(1)}(\tau)|^2 \). As is well known from the Wiener–Khintchine theorem, the Fourier transform of \( g^{(1)}(\tau) \) yields the power spectral density of the emission. In the time domain it leads to Rabi oscillations.

\[
g^{(1)}(\tau) = A_0 + A_1 e^{-i\Omega \tau} + A_2 e^{i\Omega \tau} + A_3 e^{-2i\Omega \tau} = \Omega \frac{g^{(2)}(\tau)}{g^{(1)}(\tau)}
\]

(3)

Here \( \Omega = (\Omega_0^2 + \delta^2)^{1/2} \), is the effective Rabi frequency at detuning \( \delta \). The second-order correlation function would then have the form:

\[
g^{(2)}(\tau) = A_0 + A_1 e^{-i\Omega \tau} + A_2 e^{i\Omega \tau} + A_3 e^{-2i\Omega \tau} + A_4 e^{2i\Omega \tau} + A_5 e^{-2i\Omega \tau} + \text{incoherent terms}
\]

(4)

Thus, the time-delayed intensity–intensity correlation function for the emission from a collection of \( N \) atoms will display coherent Rabi oscillations that decay at rates indicative of the relaxation mechanisms [19].

Indeed, our TDII measurements of the fluorescent emission from the cold atoms do exhibit such damped oscillations (figure 2(b)); for large detunings, five full oscillations are seen, while for small detunings barely one or two are seen. A decaying oscillatory function was fitted to the experimental data (figure 3(a)) to determine the effective Rabi frequencies at various detunings. A comparison (figure 3(b)) of the values so obtained with those theoretically expected shows the former to be lower by \( \sim 13\% \). Fitting with two or more decay
constants (as in equation (4)) and a better timing resolution in the experiment are likely to improve the agreement. The reduced prominence of the oscillations at smaller detunings may be understood in terms of the temperature of the atoms. Cooling beams closer to resonance result in a hotter collection of atoms in which the increased inter-atomic collisions decohere the system faster.

For the collection of cold atoms in our experiments, signatures of the coherent processes in the $g^{(2)}(\tau)$ die out within delays of $\sim 150$ ns. Spontaneous emission is expected to cause atoms to undergo transitions to the ground state within a duration of a few lifetimes of the excited state. For delay times larger than this, $g^{(2)}(\tau)$ is dominated by effects due to the scattering by moving atoms. While the collection of atoms under study is cooled to $\sim 100$ $\mu K$, where the Doppler width reduces below the natural linewidth, it may seem surprising that the effect of the velocity distribution is seen in the scattering. Once again, the power of the second-order correlation becomes evident, as $g^{(2)}(\tau)$ may be interpreted as the measure of the probability of detecting a second photon scattered by an atom with velocity $v$, within a time $\tau$ of having detected one such photon. When the velocity spread of atoms is large, as at higher temperatures, the probability for such an event is low, and thus $g^{(2)}(\tau)$ will fall more rapidly towards unity compared to the case for lower temperatures. Thus, $g^{(2)}(\tau)$ at large time delays (in this case delays larger than $\sim 500$ ns) may be used to determine the temperature of the ensemble. The elastically scattered light has a Doppler profile determined by the velocity distribution of atoms [14, 20]. In the six-beam configuration, denoting by $\alpha_j (=2(1 - \cos \theta_j))$ the dependence of the Doppler spread of the $j$th beam on its scattering angle $\theta_j$, and by $A_j$ the weight factor for the $j$th beam appropriate for its intensity and its polarization and angle dependent elastic scattering cross-section [14, 20],

$$g_{ij}^{(1)}(\tau) = \sum_{j=1}^{6} A_j \exp \left( -\alpha_j \omega_0^2 k_B T \tau^2 / 2mc^2 \right)$$

(5)

Here $g_{ij}^{(1)}(\tau)$ represents the Doppler contribution to the first-order correlation function and $\omega_0$, $c$, $k_B$, $T$ and $m$ are the frequency and speed of light, the Boltzmann constant, the temperature of the ensemble and the mass of the atom, respectively. Using this in conjunction with the relation

$$g^{(2)}(\tau) = 1 + S \left| g^{(1)}(\tau) \right|^2$$

(6)

$\approx 1 + S \left| g_{ij}^{(1)}(\tau) \right|^2$ for large $\tau$

(7)

where $S$ depends on the spatial coherence of the light detection system, the temperature $T$ may be estimated from the experimentally obtained time-delayed intensity correlation function. Figure 4(a) displays the experimentally obtained value for $g^{(2)}(\tau)$ as a function of $\tau$ for different detunings $\delta$ of the cooling beam. The temperature of the ensemble is determined for each detuning of the cooling beam by fitting the $g^{(2)}(\tau)$ curve obtained for different detunings, using equations (5) and (7), to the corresponding experimental data.

As seen in Figure, the data for the different detunings fit quite well with the respective curves. It may be noted that the same parameters ($S$, $\alpha$, $A_i$) are used for all curves. The values of temperature thus obtained ($T_{TDII}$), on comparison with the temperature obtained by the trap oscillation method ($T_{TO}$), show fairly good agreement (figure 4(b)), with $T_{TDII}$ being slightly lower in all cases.

We now turn our attention to the remaining observation—the variation of $g^{(2)}(0)$ with detuning. Figure 2(c) shows that while $g^{(2)}(0) = 2$ for large detuning of the cooling beams ($\sim 22$ MHz), for the case of the cooling beams closer to resonance, the value of $g^{(2)}(0)$ drops. We attribute this to the timing resolution in our experiment. The larger the detuning, the colder the collection of atoms, and hence the slower the decay in coherence. Thus, the timing resolution of 5 ns appears adequate for large detunings. For hotter atoms obtained at lower detunings, the averaging effect of the time bin becomes discernible, as it is now a larger fraction of the coherence time.

The various results presented so far clearly indicate that TDII measurements of emission from a large collection of independent atoms enables the study of their coherent and incoherent dynamics, and allows the determination of the relative strengths of various relaxation mechanisms as functions of different physical parameters. In an earlier study, single atom dynamics were probed by photon–photon correlation in an optical dipole trap [18], where one, two, or three atoms were held trapped. In the present experiment, the number of atoms contributing to the collected light is three orders of magnitude higher. Further, in the present experiment, atoms move in and out of the region from which light is collected, due to the thermal motion, and the superimposed trap oscillation in the quadrupolar magnetic field. The transit time of an atom (in the absence of a collision that expels it from this region), is estimated to be $\sim 10$ $\mu s$.

In summary, this paper presents a TDII study of light emitted by ultracold atoms in a steady-state MOT. Though illuminated by coherent light, the atoms alter the nature of light by imprinting, upon emission, signatures of the underlying physical processes. The zero-delay two-time correlation function shows bunching, with the ideal value of 2. Well-defined Rabi oscillations are seen at small time delays ($<150$ ns). These progressively diminish in amplitude to yield an exponential decay at larger time delays. TDII measurement of the second-order correlation function is a powerful technique, which reveals information about the system that is not observable in direct measurements. The technique readily provides for the study of coherent and incoherent dynamics of the atomic systems, giving a relative measure of the various dynamical processes occurring at different time scales, even from a thermal ensemble of a large number of independent atoms.

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References

[1] Hanbury Brown R and Twiss R Q 1956 Nature 177 27
[2] Baym G 1998 Acta Phys. Pol. 29 1839
[3] Lednicky R 2007 Braz. J. Phys. 37 939
[4] Foellmi C 2009 Asteron. Astrophy 507 1719–27
[5] Assmann M, Veit F, Tempel J, Bersternann T, Stolz H, van der Poel M, Hvam J M and Bayer M 2010 Opt. Exp. 18 20229
[6] LS Instruments Dynamic Light Scattering www.ls-instruments.ch/technology/dynamic/light/scattering/dls.
[7] Philips D T, Kleiman H and Davis S P 1966 Phys. Rev. 153 113
[8] Martienssen W and Spiller E 1964 Am. J. Phys. 32 919
[9] Martin A, Alibert O, Flesch J C, Samuel J, Sinha S, Tanzilli S and Kastberg A 2012 Europhys. Lett. 97 10003
[10] Satapathy N, Pandey D, Mehta P, Sinha S, Samuel J and Ramachandran H 2012 Europhys. Lett. 97 50011
[11] Satapathy N, Pandey D, Banerjee S and Ramachandran H 2013 J. Opt. Soc. Am. A 30 910
[12] Pandey D, Satapathy N, Suryabrahmam B, Ivan J S and Ramachandran H 2014 Eur. Phys. J. Plus 129 115
[13] Jurczak C, Desruelle B, Sengstock K, Courtioux J-Y, Westbrook C I and Aspect A 1996 Phys. Rev. Lett. 77 1727
[14] Bald S, Hoffmann D, Simon J and Walker T 1996 Phys. Rev. A 53 3469
[15] Nakayama K, Yoshikawa Y, Matsumoto H, Torigi1 Y and Kug T 2010 Opt. Exp. 18 6604
[16] Girish B S, Pandey D and Ramachandran H 2016 in preparation
[17] Kohns P, Buch P, Suptitz W, Csambal C and Erter W 1993 Europhys. Lett. 22 517
[18] Gomer V, Strauch F, Ueberholz B, Knappe S and Meschede D 1998 Phys. Rev. A 58 R1657
[19] Agarwal G S 2012 Quantum Optics (Cambridge: Cambridge University Press) ch 13
[20] Westbrook C I, Watts R N, Tanner C E, Rolston S L, Phillips W D and Lett P D 1990 Phys. Rev. Lett. 65 33

Figure 4. (a) $g^2(\tau)$ versus $\tau$ for different detunings, $\delta (1 = -14.5$ MHz, $2 = -16.1$ MHz, $3 = -18$ MHz, $4 = -19.4$ MHz, $5 = -22$ MHz) of the cooling beam from the cooling transition. The values obtained from TDII measurements are shown as dots. The solid curves represent the values of $g^2(\tau)$ obtained from equations (5) and (7), with temperatures 100 $\mu$K, 80 $\mu$K, 60 $\mu$K, 50 $\mu$K, 43 $\mu$K, for the five values of detunings (1 = -14.5 MHz, 2 = -16.1 MHz, 3 = -18 MHz, 4 = -19.4 MHz, 5 = -22 MHz), respectively. The data for all sets are fitted with the same choice of parameters, to obtain the temperature. (b) Comparison of the temperature of the cold atoms, obtained from two methods—TDII (circles) and trap oscillation (squares)—as a function of the (red) detuning of the cooling beam.