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The Temperature of Optical Molasses for Two Different Atomic Angular Momenta.

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Abstract. – We have measured the temperature of laser-cooled Rb atoms in optical molasses as a function of laser intensity and detuning. For both ^{85}Rb and ^{87}Rb, cooled on the $F = 3 \rightarrow F' = 4$ and $F = 2 \rightarrow F' = 3$ transitions, respectively, the temperatures are proportional to the ratio of laser power and detuning for a wide range of these parameters. We observe a small but significant difference between the two isotopes. We also show the results of three-dimensional semi-classical numerical calculations. Our results favor a model which includes atomic localization in optical standing waves.

Since the discovery of laser cooling below the Doppler limit [1], considerable theoretical and experimental effort [2-4] has gone into understanding the new mechanisms of laser cooling. These mechanisms rely on the presence of degenerate ground states and differential light shifts and optical pumping between these states. Until now, it has been an open question how the temperature, the main parameter used in the characterization of optical molasses, depends on the degree of degeneracy, determined by the ground-state angular momentum. It is, therefore, of great interest to perform an experiment in which the role of the angular momentum can be clearly isolated. The two naturally abundant isotopes ^{85}Rb and ^{87}Rb (nuclear spins 5/2 and 3/2, respectively) with the $5S_{1/2}$-$5P_{3/2}$ transition at 780 nm offer an attractive system for separating the effects of angular momentum from other parameters such as transition linewidth or wavelength. For each isotope we have measured the dependence of the temperature of three-dimensional (3D) optical molasses on laser parameters with the highest accuracy to date. Since the only action required to shift between the isotopes is a slight change in laser frequency, our comparison of cooling on the $F = 3 \rightarrow$

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$\rightarrow F' = 4$ transition of ^{85}Rb with that on the $F = 2 \rightarrow F' = 3$ transition of ^{87}Rb should be particularly immune to systematic errors. Contrary to, *e.g.*, the decay rate from the molasses, the temperature is also rather insensitive to perturbing effects such as collisions, which have little effect on the cooling process.

Of course, laser power and detuning also determine the temperature T. Within certain limits, which are fulfilled in our experiment, T is proportional to the light shift (AC Stark shift) which contains all dependence on the laser parameters, and the atomic angular momentum enters through the constant of proportionality. These constants are the subject of this paper. The simple dependence on laser parameters (the light shift is proportional to the ratio of laser intensity and detuning) has been predicted in one-dimensional calculations [5-7] and has been seen in a three-dimensional experiment on Cs [8] and in our measurements. For it to hold, the calculations require low saturation of the cooling transition and large detuning, and in the experiment we saw deviations from the simple law when these limits were violated.

To apply our results to other atoms, one needs to know how the atomic mass m and the transition wavelength λ affect the temperature. From the one-dimensional calculations one infers that for laser parameters which keep the excited-state population small (but proportionality between temperature and light shift not required), all dependence on m and λ vanishes if temperature and light shift are scaled by the recoil energy $E_{\mathrm{rec}} = (\hbar k)^2/2m$, with $k = 2\pi/\lambda$. In a regime where, as in our experiment, the temperature is proportional to the light shift, the constant of proportionality is of course independent of such scaling.

Only recently, (numerical) calculations of sub-Doppler cooling have been developed for three dimensions [9, 10]. We extended the work of ref. [9] to $F = 2$ and $F = 3$ ground states and present the results in this paper, allowing us to make the first direct, quantitative comparison between theory and experiment. We find that the procedure adopted in [9] for obtaining temperatures leads to very poor agreement with the experiments. A significant improvement, however, results from taking into account, in a simple model, the localization of the atoms.

The optical molasses in our experiment consists of three mutually orthogonal pairs of laser beams, the two beams of each pair having orthogonal linear (lin $\perp$ lin) polarizations. A titanium:sapphire laser, locked with an uncertainty of $\pm$ 500 kHz to a saturated absorption resonance in a Rb vapor cell, serves as a light source for the molasses. An acousto-optic modulator allows precise frequency shifting and fast turn-off. A beam from a laser diode tuned to the $F = 2 \rightarrow F' = 3$ (for ^{85}Rb) or $F = 1 \rightarrow F' = 2$ (for ^{87}Rb) is superposed on one or all of the molasses beams for repumping from the wrong hyperfine state (see [11] for a level scheme). We determine the temperature of the atoms with a time-of-flight (TOF) method [1]: the atoms are dropped by shutting off the molasses beams. We record their fluorescence as they fall through a 1.1 mm thick horizontal light sheet located some 4 cm below the molasses center. The size of the atomic cloud as it passes through this probe beam determines the width of the TOF signal. It depends not only on the spread of the atoms due to their initial velocity but also on the initial extension of the cloud.

Two different methods were used to reduce the contribution of the initial geometry to the width of the TOF signal. In the first set-up, similar to the one of ref. [8], the slow atoms from a diode-laser-chirp–cooled atomic beam [12] drifted directly into the molasses, filling it to its full size (18 mm diameter). After 1 s of loading at high intensity and 4Γ detuning, the laser parameters under investigation were applied for 300 ms, which is large compared to the equilibration time of less than 1 ms. Then we shut off the molasses light (transition time less than 1 μs) and for 1 ms pulsed on an auxiliary beam (slicer) with a shadow in its center which blew away all the atoms except those in a horizontal slice 3.2 mm thick. The slicer was derived from the Ti:sapphire laser and tuned to resonance using an acousto-optic modulator.

The slicer intensity had to be kept low enough to avoid heating effects due to light scattered into the dark region from the illuminated atoms. We inferred a molasses density of $\sim 10^8\,\mathrm{cm}^{-3}$ from the 10% attenuation of a resonant weak probe beam.

In the second set-up, we first loaded the atoms from the cooled atomic beam into a magneto-optic trap (MOT) [13] where they were confined to a small volume (FWHM of density profile: 0.8 mm). After a loading phase of 0.5 s to 2 s, we turned the magnetic field off (time constant 100 μs) and changed the polarization of the molasses beams with an electronically controllable liquid-crystal retarder from circular (required for the MOT) to orthogonal linear. After 100 ms, we dropped the atoms from the molasses. This period of 100 ms allowed the light polarization to stabilize (rise time of the liquid crystal approximately 30 ms), while the atoms diffuse less than 0.2 mm, according to measurements made by a video camera. The measured intensity profiles of the molasses beams were flat to ± 3% over the central 6 mm diameter. This set-up, compared to the one using the slicer, has the advantage that the atoms are confined to the centers of the laser beams and see only the uniform part of the beam. On the other hand, since the MOT can load to much higher density than the molasses, we needed to limit the MOT loading time to avoid a density effect on the temperature [14]. Because of the smaller MOT volume this generally resulted in a smaller signal than in the slicer set-up.

We fit a Gaussian to the TOF signal to obtain its FWHM. We saw no deviations, above our noise, from a Gaussian shape, except in a narrow range of laser intensities below those which gave the lowest temperatures and just above those where the molasses stopped working. When we used the slicing beam to reduce the geometric contribution to the TOF signal, a small calculated correction to the temperature was still necessary. It never exceeded 1 μK.

Figure 1 shows the measured temperatures for ^{85}Rb as a function of $\Omega^2/|\delta|\Gamma$, where Ω denotes the Rabi frequency per laser beam for the strongest transition ($\Omega^2/\Gamma^2 =$

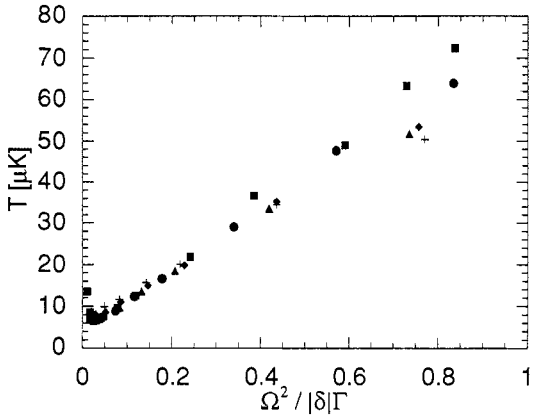

Fig. 1.

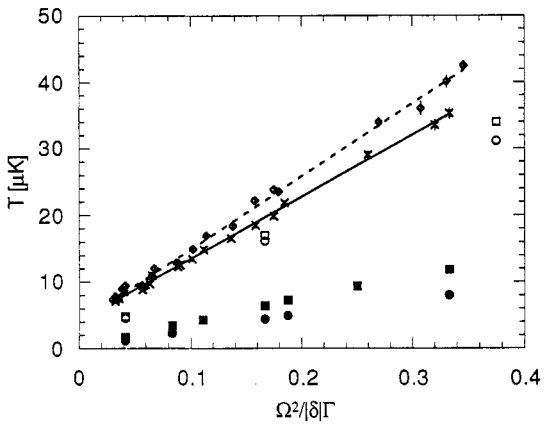

Fig. 2.

Fig. 1. – The temperature of ^{85}Rb in molasses with orthogonal linear polarization in each beam as a function of $\Omega^2/|\delta|\Gamma$ for various detunings (to the red of the atomic resonance). + 2Γ, ◆ 4Γ, ▲ 6Γ, ● 8Γ, ■ 10Γ.

Fig. 2. – Measured and calculated temperatures *vs.* $\Omega^2/|\delta|\Gamma$ for ^{85}Rb (—⋆—) and ^{87}Rb (- - -◇- - -) in the linear region with fitted straight lines. The experimental error bars are the standard deviation of the mean of ten consecutive measurements. Calculated temperatures including localization effects are represented with open squares (T_{loc}(2-3)) ($F = 2 \rightarrow F' = 3$ transition) and circles (T_{loc}(3-4)) ($F = 3 \rightarrow$ $\rightarrow F' = 4$ transition), those neglecting these effects with the same full symbols.

= *intensity per beam*/3.2 mW/cm^2), δ the laser detuning from the atomic resonance, and $\Gamma/2\pi = 5.89$ MHz, the linewidth of the atomic transition. As expected [5], the data points for different detunings in fig. 1 do indeed lie on the same curve, except those at the smallest detuning. Also, as predicted by 1D theories [5, 7], the temperature seems to depend on the detuning at large light shifts. The minimum observed temperature was 6 μK for both isotopes. If one reduces the intensity below the one giving the lowest temperature, the TOF signal decreases rapidly, and the temperature rises as expected [5, 7]. Over the parameter range where the temperature depends linearly on $\Omega^2/|\delta|\Gamma$, we define $c_\perp$ (the subscript indicating the light polarization):

$$k_B T/\hbar\Gamma = c_\perp \Omega^2/|\delta|\Gamma + \text{const}.$$

Note that, according to the 1D theories, $c_\perp$ does not depend on atomic properties other than F.

We made four independent comparisons of $c_\perp$ for ^{85}Rb and ^{87}Rb. Switching isotopes merely meant changing the locking points of the lasers; no optical realignment was done within any of the four sets. The first set was taken with the slicer set-up, and the subsequent ones used the MOT. Between these two groups of measurements, most of the optical set-up was changed. Weighted fits to each of the sets gave $c_\perp$ values of 0.33, 0.37, 0.37 and 0.33 for ^{85}Rb and 0.39, 0.41, 0.39 and 0.37 for ^{87}Rb. Our best estimate of the 1σ random and systematic uncertainty is 8%. Known systematic errors include power meter calibration, absorption in optical components, inhomogeneity of the molasses beams, and the laser frequency lock offset. For the ratios $c_\perp(^{85}\text{Rb})/c_\perp(^{87}\text{Rb})$ we obtain 0.842, 0.865, 0.936 and 0.905, giving an average of 0.89 ± 0.04, where the uncertainty is the standard deviation of the four measurements. All known systematic errors cancel for the ratio, since essentially nothing was changed when the isotopes were switched. In fig. 2 we show one data set with $0.03 < \Omega^2/|\delta|\Gamma < 0.4$ and $4\Gamma \leqslant |\delta| \leqslant 10\Gamma$, for both isotopes, together with fitted straight lines.

Our results are remarkably similar to the $c_\perp = 0.35$ found in Cs [8]. This is true not only for the slope, but also for the minimum temperature. According to the 1D theories, $T_{\min}$ for different atoms should scale as the recoil energy E_{rec}. We found $k_B T_{\min} = 16E_{\text{rec}}$, close to the value of $13E_{\text{rec}}$ for Cs [8] and $11E_{\text{rec}}$ for Na [15]. Temperature *vs.* light shift has also been measured in Na [3, 15], however, the excited-state hyperfine structure restricts the detuning to less than 3 linewidths so these data are not in the large detuning limit.

In our numerical calculations, we solve the optical Bloch equations for atoms at rest at different positions in the 3D laser field. We determine the two-time correlation function of the force operator which then allows us to calculate D_p and α. Here D_p is a position-dependent momentum diffusion coefficient which describes heating, and α is a low-velocity position-dependent friction coefficient. The equilibrium between damping and heating is characterized by a temperature $T = \overline{D}_p/k_B\overline{\alpha}$, where the position-averaged diffusion and friction coefficients both become scalar quantities under the symmetric conditions studied. There are light-field configurations where the relative phase of the light field may be adjusted so that there is no gradient of ellipticity, thus partially removing the Sisyphus cooling mechanism [9]. However, for the lin $\perp$ lin configuration this situation does not occur, so that T is not very sensitive to the phase, and we make all calculations with only one setting of phases. The details of this calculation have been described elsewhere for a $1 \rightarrow 2$ transition [9]. Here we present the extension to higher angular momenta, allowing us to calculate the temperatures for $2 \rightarrow 3$ and $3 \rightarrow 4$ transitions as measured in the experiment. Note that in this semi-classical formulation the temperature depends only on the angular momentum of the transition. The extension is straightforward, though computer-time consuming.

Our numerical 3D calculations give values of $c_\perp = 0.087$ for ^{85}Rb and $c_\perp = 0.126$ for ^{87}Rb,

resulting in a ratio of 0.69 (see fig. 2). Comparing this with the experiment, one finds the measured values of $c_\perp$ almost a factor of 4 larger, and the experiment also disagrees with the calculated ratio. Since the observed temperatures correspond to kinetic energies smaller than the light shift, we believe localization effects to be important [16-18], and we modified the calculations to take such effects into account. We no longer calculate the diffusion and friction coefficients as spatially uniform averages. Instead, we assume the spatial density of the atoms to be given by a Boltzmann distribution in a potential $\mathcal{U}$ for which we choose (see below) the expectation value $\langle V_{\mathrm{al}} \rangle$ of the laser-atom interaction, and $\bar{\alpha}$ and $\bar{D}_{\mathrm{p}}$ are now weighted averages. Since the Boltzmann distribution depends on $k_{\mathrm{B}} T = \bar{D}_{\mathrm{p}} / \bar{\alpha}$, this requires an iterative procedure. For our 3D molasses, we have calculated temperatures at $\delta = -6\Gamma$ and $\Omega^2 / |\delta| \Gamma = 0.04$, 0.17 and 0.38, from which we obtained $c_\perp = 0.31$ (^{87}Rb) and $c_\perp = 0.28$ (^{85}Rb), hence a ratio of 0.91. Note that our approximate procedure predicts values of $c_\perp$ which are in substantially better agreement with the data. Taking the localization into account increases the temperature: the weighted mean friction is reduced substantially because the number of atoms is lower in the field nodes where α diverges [9], whereas applying the weighting changes $\bar{D}_{\mathrm{p}}$ less dramatically.

This treatment of localization should be understood only as a model, and our choice of the potential $\mathcal{U}$ is somewhat arbitrary. The following discussion gives a brief summary; details will be defered to a future publication. The position-dependent mean force which localizes the atoms does not in general derive from a potential [19,20]. However, $F \equiv -\langle \nabla V_{\mathrm{al}} \rangle = -\nabla \langle V_{\mathrm{al}} \rangle / 2$ holds exactly for a number of particular field configurations (weak fields), and to a rather good approximation in more general cases that we have studied numerically. This suggests using $U = \langle V_{\mathrm{al}} \rangle / 2$ in our Boltzmann distribution, but there is another effect which influences localization as well: the position dependence of α and D_{p}. By numerically solving a realistic 1D Fokker-Planck equation, we observe that even without taking the position-dependent forces on the atoms into account, the atoms accumulate at positions which coincide with the minima of $\mathcal{U}$. The spatial correlations between $\langle V_{\mathrm{al}} \rangle$, D_{p} and α which lead to that result are not restricted to one dimension. We believe all this justifies, in a heuristic model, to account for the additional localization effect by a prefactor on the potential in the Boltzmann distribution. As it turns out, due to the simple scaling of molasses temperature with the field parameters, with this prefactor one can fit any experimental set of data points. The same factor of 2, applied here, however, yields improved agreement for both isotopes, and has also proven to be a good choice for other configurations, a discussion of which is beyond the scope of this paper.

Let us finally relate our work to some recent studies, with emphasis on different aspects or approaches to laser cooling. It has been observed that atoms may oscillate in internal-state–dependent potentials, instead of having their motion governed by (density matrix) averaged forces [17,18]. Also, full quantum calculations have revealed that, due to this kind of entangling of internal and external atomic evolution, the assumptions for a semi-classical treatment are more difficult to achieve [5,7]. Our phenomenological factor, included in the Boltzmann distribution, might also partially represent the modified localization effects in individual potentials as compared to those in a mean potential. These are ideas of more speculative character; the main point we wish to make here is that we have found that including localization in the semi-classical calculations results in a large correction to the values obtained in [9], and thus represents an important new element in the semi-classical theory.

We note also that our calculation is not able to determine T_{min}. It predicts $T \to 0$ as $\Omega_{\mathrm{Rabi}} \to 0$. In a more complete semi-classical theory one would determine the atomic momentum distribution from a Fokker-Planck equation retaining the full position and velocity dependence of both the force and the diffusion tensor. Such an approach has been initiated in 3D in ref. [10]. It would be

of interest to examine the results of this treatment for the parameters studied here as a test of the validity of our simple averaging procedure. For practical calculations, however, a complete semi-classical approach may well be so involved that, in view of recent progress [6, 21], the full quantum treatment becomes competitive.

In summary, for ^{85}Rb and ^{87}Rb in optical molasses in a 3D lin $\perp$ lin laser beam configuration we have measured the dependence of the temperature on laser parameters with the highest accuracy to date and made the first test of a model of 3D laser cooling. A theoretical treatment which neglects atomic localization disagreed with the experiment, whereas an *ad hoc* model including this phenomenon reproduced the experimental data fairly well. The high accuracy of our data should provide a rigorous test of future theoretical treatments.

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REFERENCES

[1] LETT P., WATTS R., WESTBROOK C., PHILLIPS W., GOULD P. and METCALF H., *Phys. Rev. Lett.*, **61** (1988) 169.
[2] DALIBARD J. and COHEN-TANNOUDJI C., *J. Opt. Soc. Am. B*, **6** (1989) 2023.
[3] UNGAR P. J., WEISS D. S., RIIS E. and CHU S., *J. Opt. Soc. Am. B*, **6** (1989) 2023; WEISS D. S., RIIS E., SHEVY Y., UNGAR P. and CHU S., *J. Opt. Soc. Am. B*, **6** (1989) 2072.
[4] COHEN-TANNOUDJI C. and PHILLIPS W., *Phys. Today* (1990), p. 33.
[5] CASTIN Y., DALIBARD J. and COHEN-TANNOUDJI C., in *Light Induced Kinetic Effects 1990*, edited by L. MOI, S. GOZZINI, C. GABBANINI, E. ARIMONDO and F. STRUMIA (ETS Editrice, Pisa) 1991.
[6] CASTIN Y. and DALIBARD J., *Europhys. Lett.*, **14** (1991) 761.
[7] CASTIN Y., Thesis, Paris (1992).
[8] SALOMON C., DALIBARD J., PHILLIPS W. D., CLAIRON A. and GUELLATI S., *Europhys. Lett.*, **12** (1990) 683.
[9] MØLMER K., *Phys. Rev. A*, **44** (1991) 5820.
[10] JAVANAINEN J., *Opt. Commun.*, **86** (1991) 475.
[11] SHEEHY B., SHANG S.-Q., WATTS R., HATAMIAN S. and METCALF H., *J. Opt. Soc. Am. B*, **6** (1989) 2023.
[12] WATTS R. and WIEMAN C., *Opt. Lett.*, **11** (1986) 291.
[13] RAAB E., PRENTISS M., CABLE A., CHU S. and PRITCHARD D., *Phys. Rev. Lett.*, **57** (1987) 2631.
[14] CLAIRON A., LAURENT PH., NADIR A., DREWSEN M., GRISON D., LOUNIS B. and SALOMON C., *The 5th European Frequency and Time Forum, Besançon, France, 1992.*
[15] LETT P. D., PHILLIPS W. D., ROLSTON S. L., TANNER C. E., WATTS R. N. and WESTBROOK C. I., *J. Opt. Soc. Am. B*, **6** (1989) 2084.
[16] WESTBROOK C. I., WATTS R. N., TANNER C. E., ROLSTON S. L., PHILLIPS W. D. and LETT P. D., *Phys. Rev. Lett.*, **65** (1990) 33.
[17] VERKERK P., LOUNIS B., SALOMON C., COHEN-TANNOUDJI C., COURTOIS J.-Y. and GRYNBERG G., *Phys. Rev. Lett.*, **68** (1992) 3861.
[18] JESSEN P., GERZ C., LETT P. D., PHILLIPS W. D., ROLSTON S. L., SPREEUW R. J. C. and WESTBROOK C. I., *Phys. Rev. Lett.*, **69** (1992) 49.
[19] KAZANTSEV A. P. and KRASNOV I. V., *J. Opt. Soc. Am. B*, **6** (1989) 2140.
[20] HEMMERICH A., SCHROPP D. and HÄNSCH T. W., *Phys. Rev. A*, **44** (1991) 1910.
[21] DALIBARD J., CASTIN Y. and MØLMER K., *Phys. Rev. Lett.*, **68** (1992) 580.