Evanescent-wave mirrors for cold atoms
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Radiation pressure
exerted by evanescent waves

Radiation pressure, that is exerted on cold rubidium atoms while bouncing on an evanescent-wave atom mirror, was directly observed. It was analysed by imaging the motion of the atoms after the bounce. The number of absorbed photons was measured for laser detunings ranging from 190 MHz to 1.4 GHz and for evanescent-wave angles from 0.9 mrad to 24 mrad above the critical angle of total internal reflection. Depending on these settings, velocity changes parallel with the mirror surface were observed, ranging from 1 to 18 cm/s. This corresponds with 2 to 31 photon recoils per atom. These results were independent of the evanescent-wave optical power.

This chapter is based on the publication

6.1 Introduction

Most experimental work on evanescent-wave mirrors so far has been concentrated on the reflective properties [4,5], i.e. the change of the atomic motion perpendicular to the surface [189]. This is dominated by the dipole force due to the strong gradient of the electric field amplitude. In this chapter, measurement of the force parallel to the surface are presented. It was mentioned already in the original proposal of an evanescent-wave mirror, Ref. [3], that there should be such a force. The propagating component of the wave vector leads to a spontaneous scattering force, "radiation pressure" [152,153]. To our knowledge, we presented the first direct observation of radiation pressure exerted by evanescent waves on cold atoms. Previously, a force parallel to the surface was observed for micrometer-sized dielectric spheres moving in an evanescent-wave [190]. The basic phenomenon of radiation pressure, the photon recoil momentum, was mentioned already in 1917 by Einstein, in his work on the quantisation of the electro-magnetic field [191]. It was first observed experimentally in 1933 by Frisch [192], by means of the deflection of an atomic beam with freely propagating light. In 1976, it was proposed by Roosen and Imbert to use also a beam deflection to probe the radiation pressure of an evanescent wave [193].

In our experiment, we observed the trajectory of a cloud of cold rubidium atoms bouncing on a horizontal evanescent-wave mirror. The radiation pressure appeared as a change in horizontal velocity during the bounce. We studied the average number of scattered photons per atom as a function of the detuning and angle of incidence of the evanescent wave. The latter varies the "steepness" of the optical potential.

It was discussed in Chap. 2, that due to its short extension at the order of the optical wavelength, \( \lambda_0 \), an evanescent-wave mirror constitutes a promising tool for loading low-dimensional optical atom traps in the vicinity of a dielectric surface [82,83,86,87]. It is this application which drives our interest in experimental control of the photon scattering of bouncing atoms.

In the following section, this scattering is discussed as a source of radiation pressure by the evanescent wave. Section 6.3 describes the actual experimental configuration and the imaging method used to observe bouncing atoms. Section 6.4 investigates the radiation pressure in dependence on the angle of incidence and the laser detuning, including a discussion of several systematic errors.

6.2 Photon scattering by bouncing atoms

In Chap. 2, the phenomenon of an evanescent wave was introduced. By total internal reflection of a laser beam, that is incident in the \( xz \)-plane, the evanescent wave was established in the horizontal \( xy \)-plane at the vacuum side of a glass surface, see Figs. 2.1 and 6.1. The wave vector of the evanescent wave, \( \mathbf{k} = (k_x, 0, i \delta) \), was found with a propagating component along the surface, \( k_x = k_0 n \sin \theta_i > k_0 \), where \( k_0 = 2\pi/\lambda_0 \) is the vacuum wavenumber, \( n \) is the refractive index, and \( \theta_i \) is the angle of incidence. The optical dipole potential of the atom mirror, \( U_{\text{dip}}(z) = U_0 \exp(-2kz) \), is realised by choosing a blue laser detuning with respect to an atomic resonance.
and using the exponentially decaying field amplitude that is due to the imaginary wave vector component perpendicular to the surface, \( \kappa(\theta) = k_0 \sqrt{n^2 \sin^2 \theta - 1} \). The decay length was defined as \( \xi(\theta) = 1/\kappa(\theta) \).

The number of scattered photons per bounce, \( N_{\text{scat}} = \Gamma \beta / \delta \hbar k \), was obtained in Eq. (2.19) by integrating the scattering rate of an atom along the vertical bouncing trajectory \( (z(t), v(t)) \). Note, that \( N_{\text{scat}} \) is independent of \( U_0 \), i.e., an atom climbs the exponential mirror potential up to the turning point, no matter what the maximum optical potential at the glass surface is. The “steepness” of the optical potential is determined by \( \kappa \). The steeper the potential, the shorter the time an atom spends in the light field and the smaller \( N_{\text{scat}} \). This behaviour is shown schematically in Fig. 6.3 for two different angles \( \theta \).

We expect that an absorbed photon gives a recoil momentum to the atom,

\[
P_{\text{rec}} = \hbar k_x \hat{x},
\]

which is directed along the propagating component of the evanescent wave. This was discussed, e.g. in Ref. [193]. Experimentally, we observed this effect by the altered horizontal velocity of atom clouds after the bounce. The spontaneous emission of photons during the scattering cycles leads also to heating of the cloud and thus to thermal expansion [194, 195]. Note, that the expression (6.1) is valid exactly only for a TE-polarised evanescent-wave. In TM polarisation, the wave is elliptically polarised and both the Poynting vector and the radiation pressure may be directed away from the propagation direction of the wave [101, 193]. However, with the angle of incidence close to the critical angle \( \theta_c \), also the TM wave is nearly linearly polarised and the expression (6.1) may be used.

In principle, \( N_{\text{scat}} \) is changed if other than optical forces are present. For example, the Van der Waals attraction, that was neglected in the derivation of Eq. (2.19), tends to “soften” the potential and thus to increase \( N_{\text{scat}} \). We investigated this numerically and found it to be below the resolution of our detection method.

### 6.3 Observation of bouncing atoms

The radiation pressure experiment was performed using the same optical configuration of the evanescent-wave mirror as with the bouncing fraction experiments described in the previous chapter. Also the laser systems were identical. In order to investigate radiation pressure as a function of detuning \( \delta \) in a range as large as possible with the present laser power of 28 mW, we used a TM-polarised evanescent wave. In the previous chapter it was verified by means of the bouncing fraction, that this polarisation yields a stronger dipole potential than a TE-polarised beam of the same power, see also Eqs. (2.5) and (2.6).

Two particular differences with the former setup were, however, (i) the specific use of an optical scheme to reproducibly adjust the evanescent-wave angle of incidence and, (ii) imaging of bouncing atoms instead of recording time-of-flight signals, see Fig. 6.1. A minor difference was that the magneto-optical trap (MOT) was operated at slightly larger height (6.6 mm) above the prism surface.
Figure 6.1: (a) Evanescent-wave mirror with fluorescence imaging. Magneto-optical trap (MOT), 6.6 mm above a prism, TM-polarised evanescent-wave beam (EW), camera facing in the y-direction (CCD), resonant fluorescence probe beam from above (FP).

(b) Confocal relay telescope for adjusting the angle of incidence $\theta_i$. The lenses L1 and L2 have equal focal length, $f = 75$ mm. The "object" spot (O) is imaged to the fixed evanescent wave-spot (S). A translation of L1 by a distance $\Delta a$ changes the angle of incidence by $\Delta \theta_i$. M is a steering mirror, F is the focal plane in the telescope.

(i) Angle adjustment.— Our intention was to probe the number of scattered photons, $N_{scat} \propto \xi(\theta_i)/\delta$, as a function of decay length and detuning. Therefore it was desirable to adjust the evanescent-wave angle $\theta_i$ in a well defined manner with preserved calibration. In particular, a displacement of the evanescent-wave spot due to the angle adjustment was not admissible. Such a displacement leads to a systematic error in our measurements, see Section 6.4.2.

The optical setup with which we adjusted the evanescent-wave angle is shown in Fig.6.1(b). The basic idea is to image an (hypothetical) object (O) in the laser beam (EW) to a fixed spot (S) at the prism surface, where the evanescent wave is established. The laser beam emerged from a single-mode optical fibre, was collimated and directed through a relay telescope to the prism. The angle of incidence, $\theta_i$, was controlled by the vertical displacement $\Delta a$ of the first telescope lens, L1. This lens directs the beam, whereas the second lens, L2, images it to S. A displacement $\Delta a$ leads to a variation in $\theta_i$, given by:

$$\Delta \theta_i = \frac{\Delta a}{nf} \cdot \tag{6.2}$$

The refractive index $n$ occurs here by Snell’s law for the beam entering the prism. Due to the $2f$ lens spacing the beam is again collimated at the evanescent-wave location, with a minimum waist of 335 $\mu$m at the spot $S$ ($1/e^2$ intensity radius). The focal length of both lenses (30 mm dia.) was $f = 75$ mm, which allowed for angles up to 25 mrad beyond $\theta_e$. When using larger lenses (40 mm dia., $f = 80$ mm), also angles up to 50 mrad were possible.
6.3 Observation of bouncing atoms

Figure 6.2: Fluorescence images of a bouncing atom cloud. The first image was taken 5 ms after releasing the atoms from the MOT. The contour of the right-angle prism (width 10 mm) and the direction of the EW laser beam are indicated in the first frame. For comparison the horizontal placement of the MOT is also indicated in the frame (vertical dashed line).

(ii) Imaging.— Compared with time-of-flight methods, the strength of imaging cold atoms lays in its potential of resolving possible horizontal motion of bouncing atom clouds. More specifically, changes in the horizontal motion are considered in this chapter. (Also the mutual alignment of the MOT and the evanescent-wave was facilitated using such images.)

Atoms that have bounced on the evanescent-wave mirror were detected by induced fluorescence from a pulsed probe beam in resonance with the $F_g = 2 \rightarrow F_e = 3$ transition of the D2 line. The probe beam had a diameter of 10 mm and was directed vertically downward. The fluorescence was recorded from the side by a digital frame-transfer CCD camera (Princeton Instruments) with a commercial objective of 50 mm focal length. The integration time was chosen between 0.1 ms and 1 ms, and was matched to the duration of the probe pulse. Each camera image consisted of 400×400 pixels, that were hardware-binned on the CCD array in groups of four pixels. The field of view was $10 \times 10$ mm$^2$ with a spatial resolution of 51 μm per pixel. With 15 μm pixel width, this corresponded to a magnification of 0.6.

A typical timing sequence of the experiment was as follows. The MOT was loaded from the background vapour during 1 s. After 4 ms of polarisation gradient cooling in optical molasses the atoms were released in the $F_g = 2$ ground state by closing a shutter in the cooling laser beams. The image capture was triggered with a variable time delay after releasing the atoms. During the entire sequence, the evanescent-wave laser was permanently on. In addition, a permanent repumping
beam counteracted optical pumping of the probed atoms to the $F_g = 1$ ground state. We observed no significant influence on the performance of the evanescent-wave mirror by the repumping light.

We measured the trajectories of bouncing atoms by taking a series of images with incremental time delays. A typical series with increments of 10 ms between the images is shown in Fig. 6.2. Our detection destroys the atom cloud, so a new sample was prepared for each image. The exposure time was 0.5 ms. Each image has been averaged over 10 shots. The image at 35 ms shows the cloud just before the average bouncing time, $t_b = 36.7$ ms, that corresponds to the fall height of 6.6 mm. In later frames we see the atom cloud bouncing up from the surface. Close to the prism, the fast vertical motion caused blurring of the image. Another cause of vertical blur is motion due to radiation pressure by the probe pulse. The horizontal motion of the clouds was not affected by the probe. We checked this by comparing with images taken with considerably shorter probe pulses of 0.1 ms duration.

### 6.4 The observation of radiation pressure

#### 6.4.1 Results

Radiation pressure in the evanescent wave was observed by analysing the horizontal motion of the clouds. From the camera images, we determined the centre-of-mass (COM) position of the clouds to about ±1 pixel accuracy. Such COM trajectories are shown in Fig. 6.3(b) for various angle settings of the evanescent-wave. We see clearly, that a steep optical potential, i.e. a small decay length, causes less radiation pressure than a shallow potential. For further quantitative investigation, in Fig. 6.4, the horizontal position was plotted vs. the time elapsed since release. We find that the horizontal motion is uniform before and after the bounce. The horizontal velocity changes suddenly during the bounce as a consequence of scattering evanescent-wave photons. The change in velocity is obtained from a linear fit.

In Fig. 6.5, it is shown how the radiation pressure depends on the laser detuning $\delta$ and on the angle of incidence $\theta$. The fitted horizontal velocity change has been expressed in units of the evanescent-wave photon recoil, $p_{\text{rec}} = \hbar k_0 n \sin \theta$, with $\hbar k_0 / M = 5.88$ mm/s and $n \sin \theta$, ranging between 1 and 1.03.

In Fig. 6.5(a), the detuning was varied from 188 - 1400 MHz, or $31 - 233 \Gamma$. Two sets of data are shown, taken for two different angles, $\theta_i = \theta_e + 0.9$ mrad and $\theta_e + 15.2$ mrad. This corresponds to a decay length of $\xi(\theta_e) = 2.8 \mu m$ and 0.67 $\mu m$, respectively. We find that the number of scattered photons is inversely proportional to $\delta$, as expected. The predictions based on Eq. (2.19) are indicated in the figure (solid lines).

In Fig. 6.5(b), the detuning was kept fixed at 44 $\Gamma$ and the angle of incidence was varied between 0.9 mrad and 24.0 mrad above the critical angle $\theta_c$. This leads to a variation of $\xi(\theta_e)$ from 2.8 $\mu m$ to 0.53 $\mu m$. Here also, we find a linear dependence on $\xi(\theta_e)$. The observed radiation pressure ranges from 2 to 31 photon recoils per atom. Note, that we separate this subtle effect from the faster vertical motion, in which atoms enter the optical potential with a momentum of $p_i \approx 63 \ p_{\text{rec}}$. 
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**Figure 6.3:** Radiation pressure as a function of mirror steepness. (a) A large evanescent-wave angle $\theta_i$ causes a steep potential, $U(z)$. The incident atom momentum, $p_i = 63 p_{\text{rec}}$ ($v_i = 37 \text{ cm/s}$), corresponds with a fall height of 6.6 mm. A bouncing atoms spends more time in a shallow potential, therefore scattering more photons. (b) Cloud trajectories of bouncing atoms observed by camera images for various angle settings. The symbols correspond with those of Fig. 6.4. The dashed arrow indicates increasing mirror steepness.

**Figure 6.4:** Horizontal motion of bouncing atom clouds. The centre of mass position is plotted vs. time since release. Bouncing occurred at 36.7 ms (vertical dashed line). (a) The evanescent-wave decay length was varied as $\xi(\theta_i)/\lambda_0 = \{2.40, 1.32, 1.01, 0.86, 0.76, 0.68\}$, from large to small velocity change. The detuning was 44 $\Gamma$ and the power was 19 mW. (b) Comparison of two values of evanescent-wave power, 19 mW (■, ●) and 10.5 mW (□, ○). The detuning was 31 $\Gamma$ and the evanescent-wave decay lengths were 1.87 $\mu$m (2.40 $\lambda_0$) and 0.67 $\mu$m (0.86 $\lambda_0$). Solid lines indicate linear fits.
Radiation pressure exerted by evanescent waves

**Figure 6.5:** Radiation pressure on bouncing atoms expressed as number of absorbed photons, \( N_{\text{scat}} \). (a) Detuning \( \delta \) varied for \( \xi = 2.8 \, \mu m \) (○) and 0.67 \( \mu m \) (●). (b) evanescent-wave decay length \( \xi \) varied for \( \delta = 44 \Gamma \). The laser power was 19 mW. The thin solid line is a linear fit through the first four data points. Theoretical predictions due to Eq. (2.19): two-level atom (thick solid lines), rubidium excited-state hyperfine structure and saturation taken into account (dashed lines).

In Fig. 6.4(b), we compare trajectories for 19(1) mW and 10.5(5) mW power in the evanescent wave. As expected from Eq. (2.19), there is no significant difference in horizontal motion. For a decay length of \( \xi = 2.8 \, \mu m \) both power settings lead to essentially the same radiation pressure, that is 25(3) scattered photons for 19 mW and 23(2) photons for 10.5 mW. The corresponding observations for 0.67 \( \mu m \) decay length were 13(2) and 11(1) photons, respectively.

In the previous chapter it was discussed how the optical power determines the effective mirror surface and thus the fraction of bouncing atoms. Here this was also visible in the horizontal width of imaged atom clouds. For a given evanescent wave power, there is an upper limit for the detuning, above which no bouncing can occur. For the data in Fig. 6.5(a), this threshold is calculated as \( \delta_{\text{th}} = 6.5 \) GHz for \( \xi = 0.67 \, \mu m \) and 8.1 GHz for \( \xi = 2.8 \, \mu m \). The difference in the threshold detuning is due to the Van der Waals interaction. With our laser power, Eq. (2.19) thus predicts for our mirror a minimal (average) number of 0.25 scattered photons per atom. A second threshold condition, for fixed detuning, is indeed given by the Van der Waals interaction, which yields a lower limit for the minimally useful decay length \( \xi \). For Fig. 6.5(b) this lower limit is calculated as \( \xi_{\text{th}} = 116 \) nm, i.e. for an angle \( \theta_{\text{th}} = \theta_c + 0.59 \) rad.
6.4 The observation of radiation pressure

6.4.2 Systematic errors and discussion

According to Eq. (2.19), the radiation pressure should be inversely proportional to both $\delta$ and $\kappa(\theta)$. As shown in Fig. 6.5, we find deviations from this expectation in our experiment, particularly in the $\kappa$-dependence. A linear fit to the data for $\xi < 1 \mu$m, extrapolates to an offset of approximately 3 photon recoils in the limit $\xi \rightarrow 0$ [thin solid line in Fig. 6.5(b)]. The vertical error bars on the data include statistical and systematic errors in the velocity determination from the cloud trajectories. In the following, several possible systematic errors are discussed, namely

(i) the geometric alignment, (ii) the evanescent-wave angle calibration and collimation, (iii) diffusely scattered light, (iv) the Van der Waals atom-surface interaction, (v) excited hyperfine state contributions to the optical potential, and (vi) saturation effects.

(i) Alignment. — Geometrical misalignments give rise to systematic errors in the radiation pressure measurements. For example, a tilt of the prism causes a horizontal velocity change even for specularly reflected atoms. We checked the prism alignment and found it tilted by 12(5) mrad from horizontal. This corresponds to an offset of 1.5(6) recoils on $N_{\text{scat}}$. In addition, the atoms were "launched" from the MOT with a small initial horizontal velocity, which we found to correspond to less than $\pm 0.4$ recoils for all our data.

From Fig. 6.4, we see that the extrapolated trajectories at the bouncing time $t_i$ do not start from the horizontal position before the bounce. We attribute this to a horizontal misalignment of the MOT with respect to the evanescent-wave spot. Obviously, there is a small displacement of the evanescent-wave at the prism surface, when adjusting $\theta_i$ by means of the lens L1, see Fig. 6.1(b). Since the finite-sized evanescent-wave mirror reflects only part of the thermally expanding atom cloud, such a displacement selects a nonzero horizontal velocity for bouncing atoms. We corrected for those alignment effects in the radiation pressure data of Fig. 6.5. For small radiation pressure values, the systematic error due to alignment is the dominant contribution to the vertical error bar.

(ii) Beam angle and collimation. — We checked the beam collimation and found it nearly diffraction-limited with a half-angle divergence of less than 1 mrad. The calibration of the critical angle setting was done by monitoring the power transmitted through the prism surface, while tuning the angle $\theta_i$ from below to above $\theta_c$, see Fig. 6.6. We determined $\theta_i - \theta_c$ within $\pm 0.2$ mrad. For the radiation pressure data (Fig. 6.5), the uncertainty in the evanescent-wave angle with respect to the critical angle is expressed by the horizontal error bars. Close to the critical angle, the decay length $\xi(\theta_i)$ diverges, and thus the error bar on $\xi$ becomes very large. Also the diffraction-limited divergence of the evanescent-wave beam becomes significant. It causes part of the optical power to propagate into the vacuum. In addition, the optical potential is governed by a whole distribution of decay lengths. Thus, the model of a simple exponential optical potential $\propto \exp(-2\kappa z)$ might not be valid and contribute to the disagreement of our data with the prediction by Eq. (2.19). For larger angles, i.e. $\xi(\theta_i) < 1 \mu$m, the effect of the beam divergence is negligible. This we could verify by numerical analysis.
(iii) **Diffuse light.**— Light from the evanescent-wave can diffusely scatter and propagate into the vacuum due to roughness of the prism surface. We presume this is the reason for the extrapolated offset of $\approx 3$ photon recoils in the radiation pressure [Fig. 6.5(b)]. A preferential light scattering in the direction of the propagating evanescent-wave component can be explained, if the power spectrum of the surface roughness is narrow compared to $1/\lambda_0$ [195]. The effect of surface roughness on bouncing atoms has previously been observed [194] as a broadening of atom clouds by the roughness of the dipole potential. In our case, we observe a change in centre-of-mass motion of the clouds due to an increase in the *spontaneous scattering force*. Such a contribution to the radiation pressure due to surface roughness vanishes in the limit of large detuning $\delta$. Thus, we find no significant offset in Fig. 6.5(a). Scattered light might also be the reason for the small difference in radiation pressure for the two distinct evanescent-wave power settings, shown in Fig. 6.4(b). Lower intensity of the diffuse light implies slightly less radiation pressure.

(iv) **Van der Waals interaction.**— As stated above, the Van der Waals interaction softens the mirror potential. This makes bouncing atoms move longer in the light field, thus enhancing photon scattering. This was investigated numerically by integrating the scattering rate along an atom’s path, including the Van der Waals contribution to the mirror potential. Even with the shortest decay parameter of 0.53 $\mu$m in the present experiment, the (average) number of scattered photons
would increase only about 0.8% compared with Eq. (2.19). This was not resolved experimentally. For example, with a detuning of 1 GHz and 2.5 mW power, an enhancement from $N_{\text{scat}} = 1.09$ to a value of 1.13 due to the Van der Waals interaction is calculated for a decay length of 370 nm ($\theta_i = \theta_c + 49.5 \text{ mrad}$).

However, this result was obtained by averaging the scattered photons over the effective mirror surface. At the edges of this surface, that is at the bouncing threshold circumference, $R_{\text{th}}(\phi)$ from Eq. (5.7), the turning point of a bouncing atom approaches the maximum of the mirror potential. Therefore the calculated number of scattered photons is large, i.e., it diverges for an atom at exactly the threshold circumference.

(v) Excited hyperfine state contributions.——In the two-level model, the scattering rate was expressed in the dipole potential as $\Gamma' = (\Gamma/h\delta)U_{\text{dip}}$. This is no longer true if we take into account the excited state manifold $F_e = \{0, 1, 2, 3\}$ of $^{87}\text{Rb}$. All levels except $F_e = 0$ contribute to the mirror potential and the scattering rate. Due to hyperfine pumping to $F_g = 1$, part of the atoms are lost, such that a lower net radiation pressure results. Nevertheless, we observed no influence of the permanently present repumping laser on the number of scattered photons, probably because it did not saturate the repumping transition, $F_g = 1 \rightarrow F_e = 2$.

Assuming a predominantly linear polarisation of the TM-polarised evanescent-wave with $\theta_i \approx \theta_c$, we can define a hyperfine correction, $\beta_{\text{HF}}$, to the number of scattered photons, $N_{\text{HF}} = \beta_{\text{HF}} N_{\text{scat}}$ (cf. Appendix A.3):

$$\beta_{\text{HF}} = \frac{\delta_{23}}{5} \sum_{m_g} \left( \frac{\sum_{F_e}^{+1} \left( \sum_{j' = -1}^{+1} \frac{d_{2,F_e}^2 \langle 2, m_g, 1, 0 | F_e, m_g \rangle \langle 2, m_g - j', 1, j'| F_e, m_g \rangle^2}{\delta_{2,F_e}} \right)^2}{\sum_{F_e} \frac{d_{2,F_e}^2 \langle 2, m_g, 1, 0 | F_e, m_g \rangle^2}{\delta_{2,F_e}}} \right) .$$

This correction averages over equally occupied ground state sublevels $m_g$. The numerator is proportional to the partial photon scattering rate which leaves the atom in the same ground level, $F_g = 2$. Note that the scattering amplitudes through different intermediate $F_e$ states are first added coherently, then squared [196]. The detuning for each level is assigned as $\delta_{2,F_e}$. The summation over $j'$ accounts for the three possible polarisations emitted in the scattering process. The denominator is proportional to the light shift, adding contributions from all excited $F_e$ levels. With an evanescent-wave detuning of $\delta_{2,3} = 44 \Gamma$, the correction results in a number of $N_{\text{HF}}$ scattered photons typically 9% lower than expected for a two-level atom.

(vi) Saturation effects.——In order to investigate the influence of saturation on the number of scattered photons, we solved the optical Bloch equations numerically for the steady-state excited state population, $\sigma_{\text{exc}}^{(\text{st})}$, see Appendix A.3. A bouncing atom encounters the evanescent wave as a light pulse with a typical duration between 3 and 10 $\mu$s. This is short compared to the natural excited state lifetime, $\tau = 26 \text{ ns}$. The steady-state assumption is thus justified.
The temporal variation of the Rabi frequency $\Omega_R(t)$ is expressed using the vertical bouncing trajectory, $v_z(t) = v_t \tanh(\kappa v_t t)$:

$$\zeta(t) = \frac{1}{\kappa} \ln(\cosh(\kappa v_t t)).$$  \hspace{1cm} (6.4)

Since the potential at the turning point is the maximum potential encountered by the atom, we have deliberately chosen the turning point as the origin, $\zeta = 0$, of a transformed height coordinate, $\zeta = z - (\ln U_0/U) / 2\kappa$. The Rabi frequency is then given as a function of $\zeta$, as $\Omega_R(\zeta) = \Omega_R(0) \exp(-\kappa \zeta)$, or as function of $t$:

$$\Omega_R(t) = \frac{1}{\cosh(\kappa v_t t)} = 2 \sqrt{\frac{\delta U}{\hbar}} \text{sech}(\kappa v_t t).$$  \hspace{1cm} (6.5)

We can thus integrate the time-dependent scattering rate, $\Gamma(t) = \Gamma \sigma^{\text{st}}(t)$, for a bouncing atom [cf. Eq. (A.16)]. With an evanescent-wave detuning of $44\Gamma$, we find approximately 7% fewer scattered photons compared with the unsaturated expression of Eq. (2.19). Note, that the bounces occur sufficiently slowly to preserve adiabaticity. In Fig. 6.5, we show predicted curves, corrected for hyperfine structure and saturation (dashed solid lines).

### 6.5 Conclusions

We have directly observed radiation pressure that was exerted on rubidium atoms while bouncing on an evanescent-wave atom mirror. We did so by analysing the bouncing trajectories. The radiation pressure was directed parallel to the propagating component of the evanescent wave, that is, parallel to the glass surface. We observed 2–31 photon recoils per atom per bounce and found the radiation pressure to be independent of the optical power in the evanescent wave, as expected from the exponential character of the evanescent wave.

The inverse proportionality to both the evanescent-wave detuning and the angle of incidence is in reasonable agreement with a simple two-level-atom calculation, using steady-state expressions in the limit of low saturation for the evanescent-wave optical potential and the photon scattering rate. The agreement improved when also the excited state hyperfine structure and saturation effects were taken into account. The measured number of photon recoils as a function of the evanescent-wave decay length indicates an offset of approximately 3 recoils in the limit of a very steep evanescent-wave potential. We assume, that this is due to light that is diffusely scattered due to roughness of the prism surface but retains a preferential forward direction parallel with the evanescent-wave propagating component.

With improved resolution, it should be possible to resolve the discrete nature of the number of photon recoils and also their magnitude, $\hbar k > \hbar k_0$ [197]. Our technique could also be used to observe quantum-electrodynamic effects for atoms in the vicinity of a surface, such as radiation pressure in the $xy$-plane but out of the $x$-direction of the propagating evanescent-wave component [101].