Heating and phase-space decompression of evanescent-wave cooled atoms by multiple photon reabsorption

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Abstract: We show that multiple reabsorption of resonance-frequency photons in a cloud of evanescent-wave cooled atoms can have a significant influence on the cooling efficiency and maximum value of the atomic phase-space density.

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1. Introduction

Recent progress in laser cooling and trapping of neutral atoms has opened up the possibility to study a great variety of phenomena inherent to matter at ultralow temperatures. Among the many methods of cooling, those making use of evanescent optical waves [1–7] are of particular interest when creating cold atomic samples in close proximity to a solid surface. Such samples may be used for fundamental studies of low-dimensional atomic gases [7–9], investigations of near-field properties of material surfaces [10–12], and also for applied techniques such as atom lithography [13, 14] and acceleration sensing [15–17]. In addition, microscopic atom traps and guides may be useful as integrated atom-optical devices for quantum information processing [18, 19].

The evanescent wave created by total internal reflection of a laser beam on a dielectric-vacuum interface can serve as a mirror for atoms, if the laser frequency is detuned to the blue from the atomic resonance. In such a repulsive potential, alkali atoms that comprise two ground state hyperfine levels can efficiently be made to lose kinetic energy in a so-called Sisyphus cooling process by letting the atoms repetitively bounce from the mirror potential under the effect of gravity. In this process, a low-power laser field is needed to repump the atoms back to the initial lower-energy hyperfine state in between the bounces. If the repumping beam is directed downwards, an additional slowing of the reflected atoms will take place. The combined effect of this geometric cooling process together with the Sisyphus mechanism is able to cool the atoms to a temperature corresponding to a few photon recoils [1–3]. Depending on the number density of the atoms, their phase-space density can reach a high value and allow the sample to be further processed towards quantum degeneracy [3, 6, 7].

A theoretical model describing the cooling of atoms on a tilted evanescent-wave mirror was introduced by J. Söding *et al*. in [1] and it has since been applied to calculate the cooling parameters in a variety of trap configurations [1–3, 20–22]. This model describes well the cooling of atoms at low densities or in small traps. However, the temperature achieved in experiments with large atomic samples turns out to be much higher than what the model predicts [3]. In this paper, we show that this discrepancy can be explained by the heating of the sample through the multiple reabsorption of resonant photons, emitted by the atoms in the repumping process.

In Section 2 of this work, we first briefly recall the theory of the evanescent-wave atom cooling and than calculate the contribution of the multiple photon reabsorption. Section 3 describes an example of the cooling of ^{133}Cs in a gravito-optical surface trap. In Section 4, we discuss the implication of the reabsorption effect on the cooling of atoms on an evanescent wave.

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2. Evanescent-wave cooling and multiple photon reabsorption

Evanescent-wave cooling of atoms has been experimentally demonstrated by using a gravito-optical surface trap [1–3]. This trap is created by two blue-detuned laser beams, one of which experiences total internal reflection on a flat dielectric-vacuum interface and the other one, being hollow, defines the horizontal extent of the trap (see Fig. 1). In general, the intensity of the evanescent wave created by total internal reflection decays exponentially above the surface with a decay length of $\Lambda = (\lambda/4\pi)(n^2\sin^2\theta - 1)^{-1/2}$. Here λ is the wavelength, n the index of refraction of the dielectric, and θ the angle of incidence of the reflected beam. Considering alkali atoms, we assume the evanescent wave to be tuned above the D_2-resonance frequency by an amount δ with respect to the lower hyperfine ground state $|g1\rangle$. The low-power repumping beam can be tuned into resonance with the transition from the upper hyperfine ground state $|g2\rangle$ to the excited state $|e\rangle$ (in experiments with Cs it is usually the $F = 4 \rightarrow F' = 4$ transition) in order to optically pump atoms that have entered $|g2\rangle$ back into the state $|g1\rangle$. If an atom in the state $|g1\rangle$ enters the repulsive evanescent wave and, near the turning point, makes a transition to the state $|g2\rangle$ (through scattering of an evanescent-wave photon), the energy lost by the atom as it moves up the potential will be larger than that it gains when leaving the field after the reflection. Recycling the process leads to a cooling of the atomic sample at the rate [1]

$$\gamma_{Sis} \approx \frac{2}{9}\frac{\delta_{hfs}}{\delta + \delta_{hfs}}\frac{1-q_e}{\tau_c}, \tag{1}$$

where δ_{hfs} is the ground-state hyperfine splitting and q_e is the mean branching ratio to the lower hyperfine ground state for the elastic scattering of a photon. The parameter τ_c denotes the mean time between incoherent reflections. In the case of gravitational confinement of the atoms on the evanescent-wave mirror, this parameter is given by

$$\tau_c = \frac{\hbar\delta}{\Gamma\Lambda mg\sin\varphi}, \tag{2}$$

where Γ is the natural width of the atomic line, m the atom's mass, g the acceleration of gravity, and φ the angle between the vertical axis and the vacuum-dielectric interface. In contrast to Ref. [1], the factor $\sin\varphi$ is in the denominator and, therefore, when φ approaches zero, the time τ_c goes to infinity, since in this case gravity can no longer return the reflected atoms back to the evanescent wave. Equation (2) is derived from the equation of conservation of energy, similar to that considered in Ref. [1], but containing also the potential energy due to gravity.

The rate of the geometric cooling, originating from the absorption of the downwards propagating repumping photons by the $|g2\rangle$-state atoms on their way up in the gravitational field, is calculated to be

$$\gamma_{geo} \approx \frac{4\pi\hbar\sin\varphi}{3\lambda q_r\sqrt{mk_BT}}\frac{1-q_e}{\tau_c}, \tag{3}$$

where q_r is the mean branching ratio to the state $|g1\rangle$ for transitions in the repumping field starting from the state $|g2\rangle$, k_B the Boltzmann's constant, and T the instantaneous temperature of the atoms. We point out that, in our calculations, the rate γ_{geo} in Eq. (3) turns out to be twice that given in equation (13) of Ref. [1].

Assuming that in each optical transition the kinetic energy of an atom increases by one recoil energy, one can derive an expression for the corresponding heating rate (see Ref. [1]):

$$\gamma_{heat} = \left(2 + \frac{1-q_e}{q_r}\right)\frac{(2\pi\hbar)^2}{3\lambda^2 mk_BT}\frac{1}{\tau_c}. \tag{4}$$

In thermal equilibrium, the heating rate is equal to the overall cooling rate,

$$\gamma_{heat} = \gamma_{Sis} + \gamma_{geo}, \tag{5}$$

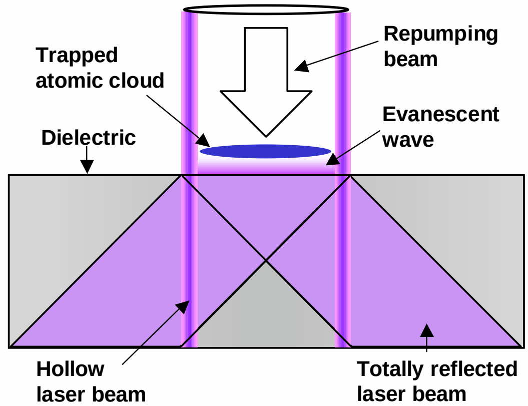

Fig. 1. A schematic view of a gravito-optical surface trap. Application of the repumping beam turns on the cooling process.

and the solution of this equation with respect to T will give the final temperature T_{eq}. Since Eq. (4) does not contain the contribution of the heating caused by the reabsorption of photons emitted by the atoms in the repumping process, T_{eq} can be considered to be correct for low atomic densities.

In order to extend the theory to include large and/or dense atomic samples, we apply the law of energy conservation and take into account multiple reabsorption events. As before, all the relevant quantities are considered in terms of their mean values per atom. We start by calculating the mean rate at which the resonant photons originally appear in the atomic cloud. In between the reflection events, the atoms spend most of the time outside the strong evanescent field and, due to repumping, they will be found in the state $|g1\rangle$ [1–3]. Thus, photons that are scattered from the evanescent wave or emitted in spontaneous transitions from the excited state to the state $|g2\rangle$ pass the sample with a negligibly small probability to be absorbed or Raman scattered, since they are far enough from resonance with the transition $|g1\rangle \rightarrow |e\rangle$. As each cooling cycle ends with a transition $|e\rangle \rightarrow |g1\rangle$, there will be resonant photons produced with a rate $R_1 = (1-q_e)/\tau_c$. The higher the density of atoms in the cloud, the higher is the probability that these photons will be absorbed by the atoms. These photons are created in spontaneous transitions and are emitted in all directions with equal probability. Since they originally appear in a very thin layer above the evanescent-wave mirror, only a fraction of them, $\eta_{in} \approx 1/2$, enter the sample. The mean rate at which the resonant photons get into the sample is, therefore, $R_{in} = \eta_{in} R_1$.

The total mean rate R_{rad} of resonant-photon emissions per atom is equal to the sum of R_{in} and R_r, where R_r is the additional rate due to multiple reabsorption events. The mean power of the resonance-frequency radiation emitted by each atom is, therefore, $P_{rad} = 2\pi\hbar\nu_1 R_{rad}$, where ν_1 is the resonant-photon frequency. A part of this radiation continuously escapes the sample at a power of $P_{out} = \eta_{out} P_{rad}$, where η_{out} is smaller than 1. Since each absorption of a resonance-frequency photon is followed by emission of another such photon, energy conservation requires that the radiation power entering the sample, $P_{in} \equiv 2\pi\hbar\nu_1 R_{in}$, has to be equal to the escaping power P_{out}. Thus we obtain

$$R_r = \left(\frac{1}{\eta_{out}} - 1\right) R_{in}. \tag{6}$$

Supposing that η_{out} is known, we now calculate the rate of heating of the atoms in the reabsorption events. This rate is given by $\gamma_r = (\Delta T_r/T) R_r$, where ΔT_r is the temperature change caused by optical recoils during the time $1/R_r$. On average, each atom in this time absorbs and

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emits one resonant photon, in total undergoing a mean number of $2/q_r$ transitions between the ground and the excited state (due to the presence of the repumping field). Each of the transitions is associated with an increase of the atomic kinetic energy by one recoil energy. Thus, $\Delta T_r = (2/q_r)(2E_{rec}/3k_B)$. By making use of Eq. (6) and writing R_{in} and E_{rec} in their explicit forms, we obtain

$$\gamma_r = \frac{1-q_e}{q_r}\frac{(2\pi\hbar)^2}{3\lambda^2 m k_B T}\frac{1}{\tau_c}2\eta_{in}\left(\frac{1}{\eta_{out}}-1\right). \tag{7}$$

The overall heating rate including the contribution of the resonant-photon reabsorption is then $\widetilde{\gamma}_{heat} = \gamma_{heat} + \gamma_r$, where γ_{heat} is given by Eq. (4). Taking into account the fact that $\eta_{in} \approx 1/2$, we obtain

$$\widetilde{\gamma}_{heat} \approx \left(2 + \frac{1-q_e}{q_r\eta_{out}}\right)\frac{(2\pi\hbar)^2}{3\lambda^2 m k_B T}\frac{1}{\tau_c}. \tag{8}$$

Compared with Eq. (4), this equation contains the parameter η_{out} in the denominator of the second term. Since η_{out} is always smaller than unity, the rate $\widetilde{\gamma}_{heat}$ is always higher than γ_{heat}, and only if resonance-photon reabsorption can be neglected is $\eta_{out} \approx 1$ and $\widetilde{\gamma}_{heat} \approx \gamma_{heat}$.

The thermal equilibrium allowing for the reabsorption effect is obtained by equating $\widetilde{\gamma}_{heat} = \gamma_{Sis} + \gamma_{geo}$. This can be written in the form showing the explicit dependence of the equilibrium temperature $\widetilde{T}_{eq}$ on the parameter η_{out}

$$\sqrt{\widetilde{T}_{eq}} = B\left\{\sqrt{1 + C\left(\frac{2q_r}{1-q_e} + \frac{1}{\eta_{out}}\right)} - 1\right\}, \tag{9}$$

where

$$B \equiv \frac{3(\delta + \delta_{hfs})\pi\hbar}{\delta_{hfs} q_r \lambda \sqrt{m k_B}} \tag{10}$$

and

$$C \equiv \frac{2q_r\delta_{hfs}}{3(\delta + \delta_{hfs})}. \tag{11}$$

The only unknown quantity in Eq. (9) is $\eta_{out} \equiv P_{out}/P_{rad}$, which may itself depend on $\widetilde{T}_{eq}$. To calculate η_{out}, we assume that the atomic sample has a constant density equal to the peak density n_0 within the effective trap volume. Then, by modelling the emitted resonance-frequency photons as spherical waves, continuously radiated by each atom at power P_{rad}, we can estimate the mean power P_{out} and thus find η_{out}. In these calculations, the attenuation of the waves with the distance of propagation has to be taken into account. The attenuation coefficient for the weak resonance-frequency radiation is given by $\alpha = \sigma_{res} n_0$, where $\sigma_{res} \approx 3\lambda^2/2\pi$ is the atomic absorption cross section corresponding to a lifetime broadened transition. The procedure for calculating η_{out} is outlined in Section 3 assuming the geometry of a horizontally aligned gravito-optical surface trap.

3. Cooling in a gravito-optical surface trap

We apply the model of Section 2 to the particular case of cooling of alkali-metal atoms in a gravito-optical surface trap. This kind of a trap is produced by separating a part of a horizontal evanescent-wave atom mirror by a vertically aligned blue-detuned hollow laser beam [2,3,23]. The atoms in the trap are confined in the vertical direction by the evanescent wave and gravity, and in the horizontal directions by the surrounding potential barrier of the hollow beam. The hollow-beam diameter D is usually on the order of 1 mm and, therefore, when the atoms are cooled down, the effective vertical size of the sample, $u = k_B T/mg$, turns out to be much smaller than D. In this case, the calculation of the parameter η_{out} becomes particularly simple, since D can be extended to infinity. The atomic sample can thus be considered to be in the form of an infinite slab with thickness u. Applying the calculation technique described in the previous Section, the average radiated power per one atom leaving the sample through the slab surfaces is obtained in the form

$$\begin{aligned} P_{out} &= \frac{2}{u}\int_0^u dz\left(\frac{P_{rad}}{4\pi}\int_0^{2\pi}\int_0^{\pi/2} e^{-\alpha z/\cos\theta}\sin\theta d\varphi d\theta\right) \\ &= P_{rad}\left(\frac{1-e^{-\alpha u}}{2\alpha u} - \frac{1}{2}[\alpha u\Gamma(0,\alpha \mathrm{u}) - \mathrm{e}^{-\alpha \mathrm{u}}]\right), \end{aligned} \tag{12}$$

where z denotes the coordinate along the normal to the slab and $\Gamma(0,\alpha\mathrm{u})$ is the incomplete gamma function. Equation (12) can in fact be cast in an approximate form, which, while still being quite accurate, allows a more transparent interpretation,

$$\frac{P_{out}}{P_{rad}} \equiv \eta_{out} \approx \frac{1-e^{-\alpha u}}{\alpha u(2-e^{-\alpha u})}. \tag{13}$$

In the case of no absorption, αu approaches zero and P_{out} approaches its maximum value of P_{rad}, as it should. On the other hand, for strong absorption, αu is much larger than 1 and P_{out} approaches zero. The attenuation coefficient $\alpha = 3\lambda^2 n_0/2\pi$ depends on the atomic density $n_0 = N/(u\pi D^2/4)$ with N being the total number of atoms in the trap. The product αu is, therefore, independent of u and Eq. (13) can be rewritten in terms of N as

$$\eta_{out} \approx \frac{1}{\beta N}\frac{1-e^{-\beta N}}{2-e^{-\beta N}}, \tag{14}$$

where $\beta = 6\lambda^2/(\pi D)^2$. We remind that this equation is valid for $u \ll D$. Substituting Eq. (14) into Eq. (9), we obtain an expression for the equilibrium temperature $\widetilde{T}_{eq}$, which will now be a function of N and D.

As an example, we consider the cooling of a sample of ^{133}Cs in a trap that was recently used in experiments by M. Hammes *et al.* [6, 7]. The inner diameter D of their hollow beam was 0.8 mm and the evanescent-wave detuning was $\delta = 2\pi \times 5$ GHz. For ^{133}Cs the ground-state hyperfine splitting is $\delta_{hfs} = 2\pi \times 9.2$ GHz, the spontaneous decay rate $\Gamma = 2\pi \times 5.3$ MHz, the resonance wavelength $\lambda = 852$ nm, the mass $m = 2.2 \times 10^{-25}$ kg, and the branching ratios are $q_e = 0.75$ and $q_r = 0.611$ [1]. Neglecting first the resonant-photon reabsorption, we calculate the equilibrium temperature of the atoms with the aid of Eq. (5) by setting $\varphi = \pi/2$, and obtain a value of $T_{eq} = 1\ \mu$K, independent of the number of the trapped atoms. The lowest temperature which has been achieved so far in the experiments with low-density atomic samples of ^{133}Cs is two times higher than this calculated value [3]. One of the possible reasons could be the ”roughness” of the realistic evanescent-wave mirror which leads to a diffuse atomic reflection [11] and to a decrease of the geometric-cooling efficiency (see $\sin\varphi$ in Eq. (3)). We now include the reabsorption processes and numerically solve Eq. (9), using Eq. (14) for η_{out}. The temperature $\widetilde{T}_{eq}$ calculated in this way is shown as a function of N in Fig. 2(a) ($\widetilde{T}_{eq}(N)$ for small values of N is shown in the inset). The resonance-photon reabsorption is seen to significantly affect the final temperature of the cooled atoms. The lowest possible temperature in cooling of, e.g., 10^8 atoms turns out to be almost two orders of magnitude higher than T_{eq} given by Eq. (5). At large values of N, which provide small η_{out}, the temperature $\widetilde{T}_{eq}$ shows a linear dependence on N with a slope given by

$$a \approx \frac{\delta+\delta_{hfs}}{\delta_{hfs}}\frac{72\hbar^2}{q_r m k_B D^2}, \tag{15}$$

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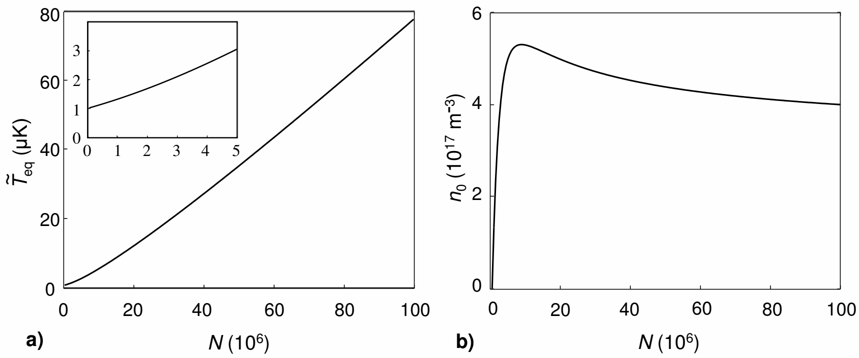

Fig. 2. Dependence of a) equilibrium temperature $\widetilde{T}_{eq}$ and b) atomic density n_0 on the number of atoms N in a gravito-optical surface trap.

as can be seen from Eqs. (9) and (14). We note that the linear dependence has been observed in experiments of Ref. [3], although the number of atoms was only $N \leq 4 \times 10^6$. If we substitute $\delta = 2\pi \times 3$ GHz and $D = 0.52$ mm as applied in Ref. [3] into Eq. (15), we obtain $a \approx 2 \times 10^{-12}$ K. This compares well with the measured value of 1.5×10^{-12} K.

The density n_0 of the atoms scales as $N/\widetilde{T}_{eq}$ ($n_0 = 4N/u\pi D^2$) and it is plotted as a function of N in Fig. 2(b). The function $n_0(N)$ has a local maximum at $N \approx 8 \times 10^6$, since at that point the relative increase of temperature, $\mathrm{d}\widetilde{T}_{eq}/\widetilde{T}_{eq}$, starts to exceed the relative growth of N. When N increases further, the density n_0 slowly drops towards a constant value. This value can be determined to be $n_\infty = 4mg/ak_B\pi D^2 \approx 3 \times 10^{17}$. However, for $N > 10^8$ the height u of the sample starts to exceed the diameter D and, therefore, Eq. (14) cannot be used anymore.

One of the important parameters characterizing the cooling efficiency is the phase-space density. For a fully unpolarized cloud of Cs, it can be expressed in terms of n_0 and $\widetilde{T}_{eq}$ as $\Omega = (n_0/7)(2\pi\hbar^2/mk_B\widetilde{T}_{eq})^{3/2}$. Figure 3a shows the dependence of Ω on N for $N < 2 \times 10^7$. The maximum value of Ω is reached at $N \approx 1.8 \times 10^6$. We note that in Ref. [6], where the parameter values are as in our calculations, the number of cooled atoms is in fact very close to this calculated optimum value ($N = 2 \times 10^6$ in Ref. [6]). Our simple model shows that Ω starts to drop fast beyond this optimum value as a consequence of the multiple photon reabsorption.

Our model indicates that it should be possible to increase the maximum value of Ω by making the effective size of the trap smaller, since then the resonant radiation would escape the sample more efficiently. In the trap considered, the height u of the sample is much smaller than the diameter D, when N corresponds to a high value of Ω. Therefore, changing D does not result in a significant change of the maximum achievable phase-space density (Ω_{max}). If, on the other hand, there would be some mechanism that would reduce the height u, the value of Ω_{max} could be increased. As an example, we show the dependence of Ω on N for atoms acted on by an additional force $\mathbf{f} = 10 \times m\mathbf{g}$ (Fig. 3b). This force can be created, e.g., by a spatially inhomogeneous far-red-detuned laser field, similar to one used in Ref. [7], but which could be applied during the cooling.

4. Discussion

We have shown that in the cooling of an atomic sample by an optical evanescent wave, the presence of photons at the resonance frequency of the atoms leads to a crucial dependence of the temperature and phase-space density of the atomic sample on the number of atoms and on the effective size of the sample. This dependence is due to the heating caused by the process

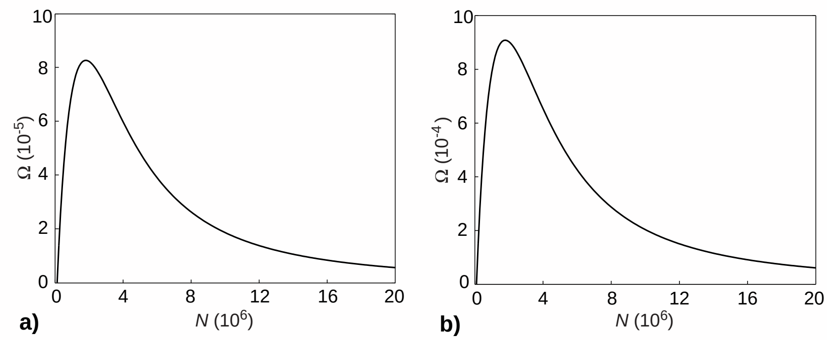

Fig. 3. Phase-space density Ω versus number of atoms N in a gravito-optical surface trap, where a) no external force is applied and b) the applied force is $\mathbf{f} = 10 \times m\mathbf{g}$.

of multiple photon reabsorption. In particular, the phase-space density versus the number of atoms in the trap shows a well defined maximum. This maximum can be increased by reducing the trap size, since the trapped resonant radiation escapes smaller samples more efficiently. Of course, the trap cannot be reduced to an arbitrarily small size for a fixed number of atoms, since at higher densities, the trap lifetime becomes shorter due to inelastic interatomic collisions. The cooling results could be improved further, if, for example, the repumping mechanism was replaced with a process that would not involve spontaneously emitted resonance-frequency photons. Such repumping could in principle be realized through Raman transitions in a field which does not contain the near-resonance frequency components.

The extension of the existing theory to include large and dense atomic samples is important, since high values of the phase-space density occur at such number densities at which the photon reabsorption has already a high rate. Although the model implies cooling of atoms on a planar surface, it can readily be adopted to more complicated traps based on blue-detuned optical fields [20, 21, 24]. The mean cooling and heating rates can in such cases be calculated by considering the atoms as a sample of ideal gas that obeys the Maxwell-Boltzmann density distribution. For traps having comparable sizes in all directions, the calculation of η_{out} will be more complicated than that given in Section 3 for the gravito-optical surface trap.

The theoretical approach described to find the reabsorption contribution can also be applied to calculate the cooling parameters for atoms confined on the evanescent wave by a far red-detuned laser beam. In this case, the effective size of the trap and, as a result, the number of atoms which could be trapped and cooled will be limited by the beam power. In addition, heating of atoms by photon recoils and loss of them due to two-body inelastic collisions (including light-induced collisions leading to molecular states) will have a higher rate. Nevertheless, in the case of a small trap size, the phase-space density can still reach a high level. To find an optimum balance between the required beam power and the detuning, multiple scattering of the photons from the red-detuned laser beam has to be taken into account in addition to the multiple reabsorption. The contribution of the multiple scattering processes to the overall heating rate can be calculated in a similar way by making use of the law of energy conservation.

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