Full length article

Loading atoms in a bi-dimensional light trap

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Abstract

We investigate theoretically a trap formed by two laser evanescent waves propagating at the surface of a dielectric prism, which confine the atoms in a Morse potential along the direction perpendicular to the prism. We consider a loading process based on the Sisyphus effect, in which a single spontaneous Raman transition is involved. We show that it is possible to achieve in this way an efficient loading of the ground state of the Morse potential, and to get thus a quasi-bi-dimensional atomic gas at the surface of the dielectric.

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1. Introduction

Using laser cooling techniques, it is now possible to prepare atomic samples cooled at the recoil limit, or even below, i.e. with a rms momentum of the order or smaller than the photon momentum $\hbar k$ [1–3]. These techniques have allowed recently a major breakthrough in the physics of quantum degenerate gas, i.e. the observation of Bose–Einstein condensation. [4–6]. One of the key elements for this realization is the possibility to confine the atomic gas without heating, which was achieved in Refs. [4–6] using a magnetic trap. This allows a further cooling of the atoms using evaporation, leading to a spectacular increase of the phase space density of the gas.

Another way for confining an atomic gas without heating is to use a laser dipole trap [7]. In this case, the confining potential is the light-shift, or AC Stark shift, of the atomic internal ground state [8]. Con-

sequently, depending on the sign of the detuning between the laser frequency and the atomic resonance frequency, the atoms are stored in locations where the light intensity is maximal [9–11] or minimal [12–14]. This detuning is in any case chosen large enough so that photon scattering processes, which cause some heating of the trapped atoms, occur only with a small rate.

In the present paper, we focus on a particular type of laser dipole trap, which provides a strong confinement of the atoms along the vertical direction z, while the atoms remain quasi-free in the horizontal plane. We show that it is possible, using Sisyphus cooling, to prepare an appreciable fraction of atoms in the ground state of the motion along the z direction. The resulting atomic sample is a 2D gas; as it is well known from statistical physics, this gas should have remarkable properties, quite different from 3D ensembles, when the quantum degeneracy factor $\xi_{2D} = \sigma \Lambda_{dB}^2$ reaches unity [15–19]. We have introduced here the surfacic density of the gas σ and the de Broglie thermal wavelength $\Lambda_{dB} = h/\bar{p}\sqrt{2\pi}$, where $\bar{p}$ is the rms momentum

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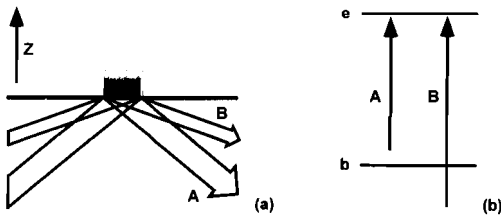
Fig. 1. Geometrical configuration studied in this paper. (a) Atoms are confined in the z direction by a Morse potential realized using two evanescent waves A and B propagating at the surface of a dielectric, with decay lengths $\kappa_A^{-1} = 2\kappa_B^{-1}$. (b) The waves A and B are detuned respectively below and above the atomic resonance transition $b-e$.

Fig. 2. Loading scheme based on a Sisyphus process. Incident atoms are prepared in state a. They are repelled by the two evanescent waves A and B which are both detuned blue from the resonance $a-e$. The atom-laser parameters are chosen such that the classical turning point is just located above the minimum of the Morse potential for state b. A Raman scattering process occurring in the vicinity of the turning point can transfer the atom to the ground state of this Morse potential.

along a given axis of the transverse motion.

The trap that we investigate is formed with two evanescent waves A and B, obtained from the total internal reflection of two incident Gaussian laser waves at the surface of a dielectric prism with index n (Fig. 1). This surface is supposed to be plane and horizontal. The wavelengths λ_A and λ_B of the two waves are of the same order and we put $k = 2\pi/\lambda_A \simeq 2\pi/\lambda_B$. The incidence angles θ_i $(i = A, B)$ are different so that the decay lengths κ_i^{-1} of the evanescent electric fields $(\kappa_i = k(n^2 \sin^2 \theta_i - 1)^{1/2})$ are also different. We choose $\theta_A < \theta_B$ so that $\kappa_A^{-1} > \kappa_B^{-1}$; in addition the wave A has a negative detuning with respect to the atomic transition of interest $b-e$, so that it attracts towards the prism the atoms in the ground state b, while the wave B has a positive detuning and repels those atoms from the prism. Consequently the total light force acting on the atoms is attractive at the long range and repulsive at the short range. There exists a plane parallel to the surface of the prism, where the force is zero, and which is located a fraction of an optical wavelength above it.

The principle of this trap using a double evanescent wave has been first proposed in Ref. [20]; in that work, the authors have analyzed classically the motion of the atoms in the trap. Here, in Section 2, we focus on the quantum aspect of the vertical motion of the atoms. We consider the particular case $\kappa_B = 2\kappa_A$, for which the potential created by the light is a Morse potential; analytical results can then be obtained concerning the number and the position of the bound states.

We then address the central point of this paper which concerns the loading of the trap. The scheme that we investigate is based on a Sisyphus effect [21,22] and
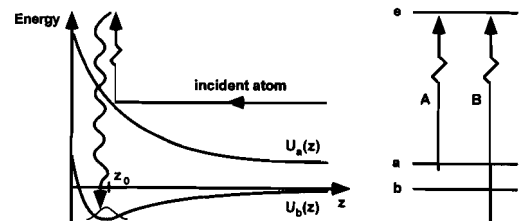
is represented in Fig. 2. We consider atoms with two internal ground states a and b, corresponding for instance to the two hyperfine states of an alkali atom. We assume that the frequencies ω_A and ω_B have been chosen such that atoms in level a see only a strong repulsive potential. On the contrary, atoms in b see the Morse potential and they can be trapped just above the surface of the prism. The loading procedure that we study theoretically works as follows. Atoms are prepared in state a, and they are dropped onto the prism on which they bounce if the repulsive potential barrier created by the lasers is larger than their initial energy. The turning point z_0 in the repulsive potential is adjusted to be at the same altitude as the minimum of the Morse potential. Around this turning point, a Raman process from a to b can occur, involving the absorption of a laser photon, and a spontaneous emission. This happens to be an efficient way for populating the ground state of the Morse potential.

We analyze this process using two different theoretical approaches. The first one is based on the Born approximation (Section 3). The second approach, presented in Section 4, consists in a numerical integration of the Schrödinger equation describing the atomic motion, the spontaneous emission processes being taken into account using the Monte Carlo wave function approach [23–25]. The two approaches are found to be in good agreement within their validity domain. They indicate that an appreciable trapping probability into the ground state of the Morse potential is achievable using this Sisyphus process.

To end this presentation, we note that other schemes

have been proposed and/or demonstrated to take advantage of the short spatial scale of an evanescent laser field and to achieve an efficient cooling process. A simple Sisyphus cooling process can occur in a single evanescent wave for three-level atoms a, b, e such as those considered here. The evanescent wave has to be blue detuned with respect to both transitions a–e and b–e, and the repulsive potential is then stronger for atoms in level b, since the detuning for the b–e transition is smaller than for the a–e transition. Atoms incident in level b may undergo a spontaneous Raman transition towards level a close to the turning point z_0; they loose in this case an appreciable fraction of their initial kinetic energy, although the final motion is still asymptotically free, contrary to the situation considered in the present paper. This process has been studied both theoretically [26,27] and experimentally [28,29]. It has also been proposed to achieve a trap with a Morse potential similar to the one considered here, and to load it directly from a magneto-optical trap located at the contact of the dielectric prism [30]. Finally, it has been suggested recently [31] to use an evanescent wave to decelerate incident atoms, which are then transferred through a spontaneous Raman transition at the antinode of a laser standing wave located in the vicinity of the dielectric, achieving therefore a quasi-bi-dimensional gas similar to the one considered here.

2. Trapping in the Morse potential

In this section, we derive the potential created by the pair of evanescent waves shown in Fig. 2 onto the atoms, depending on their internal state a or b. For state b, this corresponds to a Morse potential, and we briefly recall the number and position of the corresponding bound states. We assume that the evanescent laser fields $E_i \exp(-\kappa_i z)$ ($i = A, B$) do not saturate the atomic transitions, i.e. the maximal Rabi frequency along an atomic trajectory $\Omega_i(z_0) = dE_i \exp(-\kappa_i z_0)/2\hbar$ is small compared to the detuning of the wave i with respect to the transitions a–e and b–e. We suppose for simplicity that the reduced atomic dipole d is the same for both transitions. This is the case for alkali atoms provided the evanescent field is linearly polarized and provided the detunings are large compared to the hyperfine excited state

structure. The action of the lasers onto the atomic motion can then be described by a potential equal to the total light-shift of the occupied level $\eta = a, b$ [8]:

$$U_\eta(z) = \sum_{i=A,B} \frac{\hbar \Omega_i^2(z)}{4(\omega_i - \omega_{\eta e})}, \tag{1}$$

where ω_i is the frequency of the wave $i = A, B$, and $\omega_{\eta e}$ is the frequency of the atomic transition $e \leftrightarrow \eta = a, b$.

We choose here

$$\omega_{ae} < \omega_A < \omega_{be} \quad \text{and} \quad \omega_{ae} < \omega_{be} < \omega_B. \tag{2}$$

Therefore both waves tend to repell the atoms in state a. We assume that the effect of the wave A is dominant for those atoms so that:

$$U_a(z) \simeq \frac{\hbar \Omega_A^2(0)}{4(\omega_A - \omega_{ae})} e^{-2\kappa_A z}. \tag{3}$$

For level b, the two waves A and B have opposite effects, A attracts the atoms towards the prism and B repels them. For $\kappa_B = 2\kappa_A$, $U_b(z)$ is a Morse potential. With a prism of index $n = 1.5$, this is achieved taking for instance $\theta_A = 48°$ and $\theta_B = 70°$; we get in this case $\kappa_B = 2\kappa_A = k$. The depth U_0 and the position of the minimum z_0 of the Morse potential are given by:

$$U_0 = \frac{\hbar \Omega_B^2(0)}{4(\omega_B - \omega_{be})} e^{-4\kappa z_0} = \frac{1}{2} \frac{\hbar \Omega_A^2(0)}{4(\omega_{be} - \omega_A)} e^{-2\kappa z_0}, \tag{4}$$

where we have put $\kappa = \kappa_A$.

The number of bound states in the Morse potential is given by the integer part of $\sqrt{mU_0/(2\hbar^2\kappa^2)} + 1/2$, where m is the atomic mass [32]. We choose in the following:

$$U_0 = 50\hbar^2\kappa^2/m, \tag{5}$$

which corresponds to 5 bound states. Each eigenstate $|n\rangle$ ($n = 0, 1, \dots$) of the Hamiltonian describing the motion in the Morse potential corresponds to a wave function $\sqrt{\kappa}\chi_n(\kappa(z - z_0))$, which can be expressed in terms of the degenerate hypergeometric function. The energies associated with these bound states are given by [32]

$$E_n = -U_0 \left(1 - \sqrt{\frac{2\hbar^2\kappa^2}{mU_0}} \left(n + \frac{1}{2}\right)\right)^2. \tag{6}$$

Using (5) and $\kappa = k/2 = 3.7 \times 10^6$ m^{-1} for cesium atoms, we find that the positions E_n/h of the five bound states are -41.6 kHz, -25.1 kHz, -12.8 kHz, -4.6 kHz and -0.5 kHz.

The choice (5) associated with relations (4) determines the two ratios $\Omega_A^2(z_0)/(\omega_{be} - \omega_A)$ and $\Omega_B^2(z_0)/(\omega_B - \omega_{be})$. The final constraint on our parameters originates from the requirement that the turning point of the trajectories in $U_a(z)$, for a given asymptotic momentum p_i, should coincide with z_0. This determines the values of $\Omega_A(z_0)$ and ω_A, in terms of the two frequency scales of the problem, $\hbar\kappa^2/m$ and $\Delta = \omega_{be} - \omega_{ae}$. For instance, for a typical incident momentum $p_i = 25\hbar\kappa$, we get:

$$\Omega_A(z_0) = 18 \left(\frac{\hbar\kappa^2}{m}\Delta\right)^{1/2}, \quad \omega_A - \omega_{ae} = \Delta/4. \quad (7)$$

For sodium and cesium atoms, this value for p_i corresponds to atoms dropped respectively from a height of 7 mm and 100 μm.

We have neglected in the present reasoning the effect of the van der Waals interaction between the atoms and the dielectric prism, varying as $1/z^3$. This approximation is valid provided the minimum of the Morse potential z_0 is far enough from the dielectric surface. This is illustrated in Fig. 3 where we have taken $\kappa z_0 = 2$ and plotted the potential for cesium atoms, with and without the van der Waals interaction [33]. We have also neglected in our analysis the effect of gravity. This is valid since the variation of the gravitational potential mgz is small compared to that of the Morse
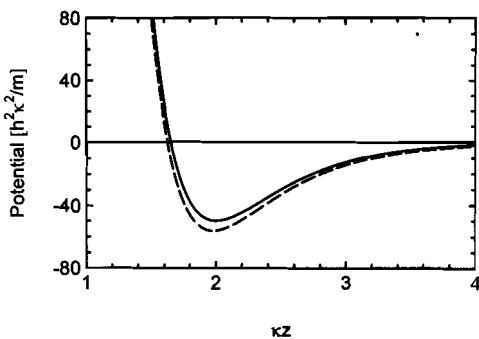
Fig. 3. Spatial variations of the potential $U_b(z)$ without (continuous line) and with (dashed line) the van der Waals interaction between the atoms and the dielectric prism. This figure corresponds to the case of cesium atoms, with a refractive index $n = 1.5$ for the dielectric material.

potential on the length scale κ^{-1}. The modification of the position of the bound states of the Morse potential can therefore be neglected. One should keep in mind however that the atoms are not completely free when they emerge from the evanescent wave region and that they will eventually slow down and turn back towards the dielectric prism.

3. The trapping probability in the Born approximation

We now turn to the calculation of the probability for an atom incident in state a, with asymptotic momentum p_i, to be transferred into a bound state $|b, n\rangle$. The incident state $|p_i\rangle$ corresponds to the eigenfunction of the Hamiltonian $H = (p^2/2m) + U_a \exp(-2\kappa z)$ [34]:

$$\frac{1}{\sqrt{L}}\Psi_i(\kappa(z - z_0))$$

$$= \left(\frac{2P_i}{\pi L}\sinh(\pi P_i)\right)^{1/2} K_{iP_i}\left(P_i e^{-\kappa(z - z_0)}\right), \quad (8)$$

where $P_i = p_i/\hbar\kappa$ and where K_{iP} is the Bessel K-function of the imaginary parameter iP. These wave functions are normalized in a box between $z = 0$ and $z = L \gg \kappa^{-1}$.

This calculation is done using the second order time dependent perturbation theory and dividing the transition rate by the incident flux $p_i/2ML$, which corresponds to the Born approximation for this problem. The transition rate for the absorption of a laser photon and the spontaneous emission of a Raman photon with wave vector k_f and polarization ϵ_f, can be written in this approximation:

$$W_n =$$

$$\frac{2\pi}{\hbar}\left|\sum_e \frac{\langle b, n; k_f, \epsilon_f|V|e\rangle \langle e|V|a, p_i\rangle}{E_a + \hbar\omega_A - E_e}\right|^2 \rho(E_f). \quad (9)$$

The operator V describes the atom–field coupling, responsible for absorption, spontaneous and stimulated emission processes; $\rho(E_f)$ corresponds to the final density of states for the fluorescence photon.

The sum (9) runs over all possible excited states e, with energy E_e. In the following, we assume that

ω_A is chosen sufficiently close to an atomic resonant frequency so that only a single excited state e contributes [1]. After summation over $\mathbf{k}_f$ and $\boldsymbol{\epsilon}_f$, Eq. (9) can then be simplified to give:

$$W_n = \Gamma_{eb} \frac{\Omega_A^2(z_0)}{4\delta_A^2} \frac{1}{\kappa L} |\mathcal{O}_n(p_i)|^2 , \qquad (10)$$

where Γ_{eb} is the decay rate from e to b, $\delta_A = \omega_A - \omega_{ae}$, and where $\mathcal{O}_n(p_i)$ describes the overlap between the incident wave function and the final one, taking into account the exponential variation of the field:

$$\mathcal{O}_n(p_i) = \int \chi_n(u) \, e^{-u} \Psi_i(u) \, du . \qquad (11)$$

Dividing the transition rate W_n by the incident flux $p_i/2mL$, we obtain the trapping probability $\mathcal{P}_n$:

$$\mathcal{P}_n = \Gamma_{eb} \frac{\Omega_A^2(z_0)}{4\delta_A^2} \tau |\mathcal{O}_n(p_i)|^2 , \qquad (12)$$

where $\tau = 2m/\kappa p_i$ is the typical *bouncing time*, i.e. the time spent by an atom with velocity p_i/m in the region of extension κ^{-1} where the evanescent field is significant. As expected, this probability does not depend of the length L of the quantization box.

The result (12) can also be written:

$$\mathcal{P}_n = n_{\text{scat}} |\mathcal{O}_n(p_i)|^2 , \qquad (13)$$

where

$$n_{\text{scat}} = \frac{\Omega_A^2(z_0)}{4\delta_A^2} \Gamma_{eb}\tau = \frac{\Gamma_{eb}}{\delta_A} P_i \qquad (14)$$

is the mean number of photons which would be scattered per bounce for a closed a–e transition and with a decay rate for e equal to Γ_{eb} [35,12]. Actually in the situation considered here, the number of photons emitted on the e–b transition by a given atom cannot exceed 1, and the result (13) is meaningful only when n_{scat} is small compared to 1. Otherwise, the atom has a large probability to decay to b before it reaches the location of the Morse potential well. In the next section, we will show how to overcome this hypothesis using the Monte Carlo wave function analysis.
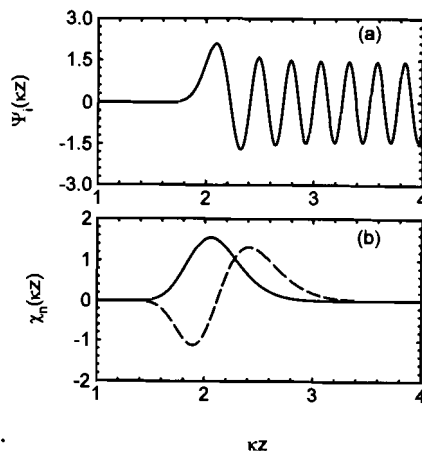
[1] For the specific case of an alkali atom, cesium for instance, this means that the detuning between ω_A and the D_2 (or D_1) resonance line is small compared to the fine structure splitting, but large compared to the hyperfine structure of the $P_{3/2}$ (or $P_{1/2}$) multiplet.


Fig. 4. (a) Spatial variations of the eigenstates of the Hamiltonian H_a corresponding to an asymptotic momentum $25\hbar\kappa$. (b) Spatial variations of the two lower states of the Morse potential χ_0 (continuous line) and χ_1 (dashed line), adjusted such that the minimum of the Morse potential coincides with the classical turning point for state a.

For a given $n_{\text{scat}} \ll 1$, we now consider the overlap factor $\mathcal{O}_n(p_i)$. In Fig. 4, we have plotted the spatial variations of the function $\Psi_i(z)$ for $p_i = 25\hbar\kappa$ (Fig. 4a). We have also plotted (Fig. 4b) the spatial variations of the two lower states of the Morse potential ($n = 0$ and 1), adjusted as explained in Section 1 so that the minimum of the Morse potential coincides with the turning point z_0. Since the first lobe of the function Ψ_i has a size similar to the extension of the ground state of the Morse potential, this leads to a large overlap factor $\mathcal{O}_0(p_i) = 0.55$. The overlap with the state χ_1 is smaller $\mathcal{O}_1(p_i) = -0.25$. The variations of the square of the overlap factor $\mathcal{O}_0(p_i)$ with the incident momentum p_i are given in Fig. 5.

In a real experiment the atoms are not necessarily prepared with a well defined momentum p_i. This causes a variation of the position of the turning point z_0, which does not coincide anymore with the bottom of the Morse potential. To estimate the corresponding reduction of the transfer efficiency, we have plotted in Fig. 6 the variations of $|\mathcal{O}_0(p)|^2$, when the incident momentum p is varied around $25\hbar\kappa$. This shows that a dispersion of a few $\hbar\kappa$, which is typical of a laser cooled atomic source, does not dramatically change the conclusions drawn above concerning the trapping probability.
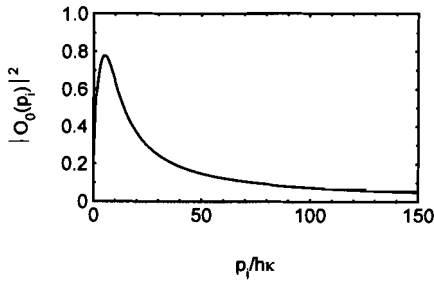
Fig. 5. Square of the overlap factor $\mathcal{O}_0(p_i)$ between the incident wave function Ψ_i corresponding to the asymptotic momentum p_i and the ground state of the Morse potential χ_0. The minimum of the Morse potential is adjusted in order to coincide with the classical turning point for p_i. The trapping probability derived from the Born approximation is equal to this quantity times the average number of scattered photons n_{scat}.
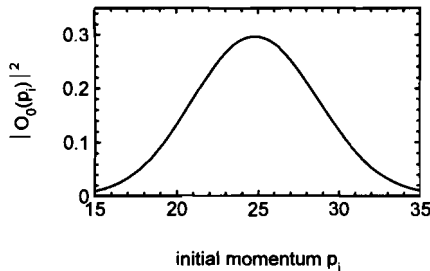

Fig. 6. Variations of the trapping probability with the incident momentum p, when the Morse potential has been adjusted such that its minimum coincides with the classical turning point for state a and for incident momentum $p_i = 25\hbar\kappa$.

4. Quantum Monte Carlo analysis of the Sisyphus process

The Born approximation presented in the previous section is valid only when the predicted number of scattered Raman photons per bounce n_{scat} (Eq. (14)) is small compared to 1. Actually, we expect that the loading of the Morse trap is optimum when n_{scat} is above 1, so that most atoms "are given a chance" to decay during their bounce. On the other hand, for an efficient loading of the χ_0 state, the decay should not occur before the atom has reached the turning point region, which imposes that n_{scat} should not be too large compared to 1. In order to determine the optimal value for this number, we have developed a Monte Carlo wave function analysis of the bouncing process [23–25].

We consider incident atoms prepared in state a far

from the surface. The initial state is a Gaussian wave packet, with an average momentum $p_i = 25\hbar\kappa$, an average position $z_i = 12.5\kappa^{-1}$ and a momentum dispersion $0.5\hbar\kappa$. This wave packet is propagated according to the Schrödinger equation with the non-Hermitian Hamiltonian:

$$H = \frac{p^2}{2m} + U_a(z) - \frac{i\hbar}{2}C^\dagger C, \qquad (15)$$

where the jump operator C is defined by

$$C = \sqrt{\Gamma_{eb}} \; \frac{\Omega_A(0)\,e^{-\kappa z}}{2\delta_A} \; |b\rangle\langle a| . \qquad (16)$$

The non-Hermitian part describes the dissipation occurring because of photon scattering processes. We consider here only Raman processes, which pump the atoms from a to b, and we neglect the atomic recoil. To be rigorous, one should also take into account the Rayleigh processes, in which the atom ends up in state a after the scattering event. However we do not expect that those processes will play a significant role, except for a slight heating due to photon recoil.

The wave function Ψ is propagated using a 4th order Runge–Kutta algorithm. At each time step, we evaluate the decrease dP of the norm of the wave function ($dP \ll 1$), which gives the probability for a scattering event. We compare dP with a random number ϵ equally distributed between 0 and 1. If no scattering occurs ($\epsilon > dP$), the wave function is just renormalized and the evolution continues. If a scattering occurs, we apply onto the wave function the jump operator C, we renormalize the result and we project it onto the states $|n\rangle$. We store the coefficients $|\langle n|\Psi\rangle|^2$, which give the probability for the atom to end up in the nth bound state of the Morse potential.

This time evolution of the wave function is illustrated in Fig. 7. As long as the atom is localized far from the evanescent wave region (Fig. 7a), the probability for a Raman scattering process is very low, because of the exponential factor $\exp(-\kappa z)$ appearing in the jump operator C. When the wave function is localized close to the classical turning point (Fig. 7b), a scattering event may occur (Fig. 7c) projecting the atom in the b internal state. If no scattering occurs, the atom is reflected back (Fig. 7d); we then stop the time evolution when the center of the wave packet reaches the starting point of the evolution.
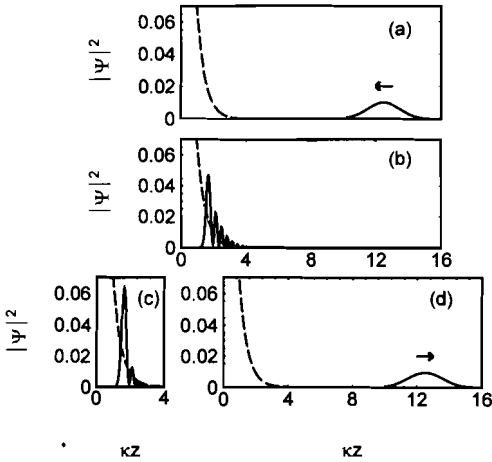
Fig. 7. Time evolution of the atomic density probability, obtained using a Monte Carlo wave function approach. (a) Incident Gaussian wave packet (internal state a), (b) wave packet around the classical turning point (internal state a), (c) wave packet after a spontaneous Raman process (internal state b), (d) outcoming wave packet if no spontaneous Raman process has occurred during the bounce. The potential $U_a(z)$ is plotted by a dashed line.
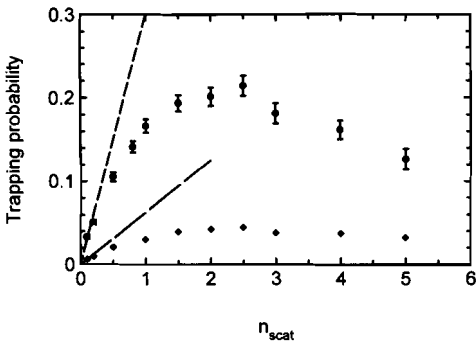

Fig. 8. Variations of the trapping probabity in state χ_0 (•) and χ_1 (◇), for an incident momentum $p_i = 25\hbar\kappa$, calculated from a Monte Carlo wave function analysis. The error bars indicate the statistical fluctuations of the results. The dashed lines represent the prediction of the approach based on the Born approximation.

This procedure is repeated 100 times, for a given value of n_{scat} (Eq. (14)). The probabilities for ending into the two lower states of the Morse potential are given in Fig. 8 as a function of this number. The error bars indicate the fluctuations of the results due to the noise inherent in the Monte Carlo method. The maximal loading probability of the ground state is 21%, obtained for $n_{scat} = 2.5$. At this value, the probabilities for ending in the low excited states of the Morse

potential are 4.5% for $n = 1$ and 0.73% for $n = 2$.

We have also indicated in Fig. 8 the prediction derived using the Born approximation: $n_{scat}|\mathcal{O}_n(p_i)|^2$ for $n = 0$ and 1. As expected, we find a good agreement between the Monte Carlo method and this prediction for low values of n_{scat}, and we confirm that the optimal value for n_{scat} is outside the range of validity of the Born approximation.

Finally we note that when $\Omega_A(z_0)$ and ω_A have been adjusted according to condition (7), the parameter n_{scat} is fixed to the value $100\Gamma_{eb}/\Delta$. For alkali atoms, this quantity is much smaller than the optimum $n_{scat} \simeq 2.5$. In practice, to circumvent this problem, one can mix with A another weak laser wave, resonant with the a–e transition, whose intensity is adjusted so that n_{scat} equals the desired value, without perturbing significantly the light-shift potential.

5. Conclusion

We have presented in this paper an efficient way for loading an atom trap which consists of two evanescent waves propagating at the surface of a dielectric and which confines the atoms via a Morse potential. We have shown that a noticeable fraction of the incident atoms (21% for $p_i = 25\hbar\kappa$) ends up into the ground state of the Morse potential. The loading probability of the excited states of the Morse potential is much smaller, which leads to a quasi-bi-dimensionnal gas.

We have restricted our treatment to the atomic motion perpendicular to the plane of the evanescent waves. The confinement in this plane can be achieved using the dipole force resulting from the Gaussian character of the laser beams forming the evanescent waves. It can also be provided using an additional quadrupole static magnetic field.

The lifetime of the trapped atoms may be limited by the off-resonant photon scattering from the evanescent fields. This scattering can be made arbitrarily low for the wave B, by choosing a large detuning $\delta_B = \omega_B - \omega_{be}$, and an accordingly large intensity Ω_B^2. One then takes advantage of the fact that the photon scattering rate varies like Ω_B^2/δ_B^2 while the confining potential varies like Ω_B^2/δ_B. For the wave A on the contrary, the detuning with respect to the transition b–e has to be chosen smaller than the splitting a–b. In the case of cesium atoms, this splitting corresponds to the

hyperfine structure of the ground level (9.2 GHz); to achieve the Morse potential studied in this paper (5 bound states), one has to take an intensity for A which leads to a photon scattering rate of ~ 500 photons per second [2]. If needed one can reduce this rate by changing rapidly the parameters of A once the loading sequence is over.

The detection of the trapped gas can be made using a pump-probe technique, comparable to the one used to study optical lattices [36]. The pump and the probe should be directed along the z axis and amplification or absorption of the probe should occur when its frequency differs from the pump frequency by an amount equal to $|E_0 - E_n|$, which is the range of a few tens of kHz for cesium atoms.

As pointed out in the introduction, this scheme provides an efficient way to obtain a degenerate bi-dimensional gas, even when starting from a non-degenerate 3D source. Indeed the surface density σ_s after loading is given by $\sigma_s = \rho H P_0$ where ρ is the volume density of the source, H its vertical size, and P_0 the trapping probability in the ground state of the Morse potential. Even if $\xi_{3D} = \rho \Lambda_{dB}^3 \ll 1$ (non-degenerate 3D gas), we can reach degeneracy for the 2D gas ($\xi_{2D} = \sigma_s \Lambda_{dB}^2 \geq 1$) if $H P_0 \xi_{3D} \geq \lambda_{dB}$. This should be feasible if ξ_{3D} is not too small, since H is a macroscopic length, much larger than λ_{dB}. In this reasoning we have omitted the fact that when an atom is loaded into the 2D gas, a photon is emitted which can be reabsorbed by another atom already present in the gas; this second atom may then be ejected from the 2D gas. This problem of reabsorption, which puts a limit onto the achievable phase space density by radiative cooling techniques, is essential in a 3D geometry [37]. For the present 2D geometry however, this problem is much less severe since most photons are emitted with an appreciable momentum along the z axis, in a direction where no atom is present.

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