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AN ALL OPTICAL TRAP BASED ON QUASIDARK STATES[1]

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We propose an all optical trap using two atomic transitions starting from a common atomic ground state with angular momentum J to two excited states with angular momenta $J_{e_1} = J_{e_2} = (J \text{ or } J-1)$. The atoms are spatially trapped and cooled towards a single quantum state which is largely decoupled from the light fields. Calculations for two blue detuned counterpropagating σ_+ and σ_- polarized light fields in one dimension predict that close to 100% of the atomic population can be trapped in the lowest quantum state. Due to the relative 'darkness' of this state spontaneous scattering of photons as well as dipole-dipole interaction between different atoms can be expected to be several orders of magnitude less than in a conventional magneto optical trap (MOT). From this one could expect extremely high densities.

1. Introduction

The past few years have seen dramatic improvements in laser cooling and trapping of free atoms [1]. A particular goal is to increase the density and, at the same time, to decrease the temperature of an atomic sample in order to observe many–particle and quantum statistical effects. In the most advanced experiments, densities which are only by a bit more than one order of magnitude off from the limit of Bose-Einstein condensation have been reached. [2]. The most commonly used device for laser cooling and trapping of neutral atom is the magneto optical trap (MOT) proposed by Dalibard and coworkers [1]. Due to its relative simplicity and robustness it has become almost a standard tool in experimental quantum optics involving gaseous atomic media. It has been found that the maximum achievable density in such a MOT is to a large extend limited by a repulsive force between different atoms mediated by spontaneously rescattered photons. One way to partly circumvent this problem in Alkali atoms was demonstrated by Ketterle with his dark spot trap, by blocking the repumping beam in

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the center of the trap [3]. This auxiliary laser (sometimes just generated as a sideband of the trapping laser) is necessary to transfer atoms, which have randomly fallen into a hyperfine level decoupled from the cooling and trapping beams, back into the coupled system. By shadowing this second beam from the trap center, a large amount of the atomic population can be accumulated in the decoupled hyperfine levels. As these atoms do not scatter photons, the total density of the trapped atoms can be significantly increased.

A conceptually different mechanism to cool atoms to extremely narrow momentum distributions, that are much smaller than a single photon momentum, was proposed and demonstrated by Aspect *et al.* [4]. This method is based on the existence of coherent superpositions of the Zeeman sublevels $m=\pm 1$ of e.g. an angular momentum $J_g=1$ to $J_e=1$ transition, which decouple from the light of two counterpropagating lasers by quantum interference. These states are, however, delocalized improper (non-normalizable) states, since we require that they are eigenstates of the kinetic part of the Hamiltonian. Hence an atom which, by optical pumping, accidentally falls into this state will remain there. Consequently, with growing interaction time between light and atoms more and more atoms will be trapped in this state of well defined kinetic energy. This method is called velocity selective coherent population trapping (VSCPT). Recently this tecnique has been generalized to two and three spatial dimensions [5]. Due to the limited interaction time between the atoms and fields, the overall efficiency of this cooling mechanism is relatively low. One of the major problems connected with this cooling scheme is, that the atoms are not confined to a small volume in space, but will slowly spread due to their plane wave components with different momenta. This feature suggests the application as a coherent atomic beam splitter [4], but prevents the system to reach a finite-sized steady state. One way to circumvent this problem is to confine the atoms in space by an additional far off resonance dipole trap (FORT) [6]. Of course this immediately implies that the spatial components of the eigenfunctions cannot be given by plane waves any more. Hence the state, which corresponds to the free dark state will acquire a finite lifetime. Nevertheless, by carefully designing the corresponding dipole trapping potential, e.g. as a large flat bottom trap, this lifetime can become very large, so that an atom on average still stays most of the time in such approximate dark states. Other possible schemes rely on crossed focussed beams in the far red detuned limit.

2. Localized quasi–dark states

In this work we shall propose a different extension of VSCPT, which will lead to very efficient trapping and cooling of atoms to spatially well localized quasi–dark states. The central idea is the simultaneous use of two atomic transitions with frequencies $\omega^{A}_{1,2}$ starting from a common ground state, which we will for simplicity choose to posses $J_g=1$. The upper states corresponding to different hyperfine structure components have both angular momentum $J_{e_1}=J_{e_2}=1$. Similiar setups using higher angular momenta are equivalently possible. Furthermore, we provide two pairs of counterpropagating

$\sigma_+-\sigma_-$–polarized laser beams. The electric field of (this 1D laser) configuration reads:

$$\begin{aligned}\vec{E}(z,t) &= \mathcal{E}_1\left(a^1_+(z)\vec{\epsilon}_+ + a^1_-(z)\vec{\epsilon}_-\right)e^{-i\omega^{\mathrm{L}}_1 t} + \mathcal{E}_2\left(a^2_+(z)\vec{\epsilon}_+ + a^2_-(z)\vec{\epsilon}_-\right)e^{-i\omega^{\mathrm{L}}_2 t}\\ a^{1,2}_\pm(z) &= e^{\mp i k_{1,2} z}\end{aligned} \tag{1}$$

$\mathcal{E}_{1,2}$ denote the electrical field amplitudes, and $\omega^{\mathrm{L}}_{1,2}$ ($k_{1,2}=\omega^{\mathrm{L}}_{1,2}/c$) are the frequencies (wavevectors) of the two pairs of laserbeams. By $\vec{\epsilon}_\pm$ we denote spherical unit vectors. We assume that each laser pair interacts only with its respective atomic transition, which enables us to eliminate the explicit time–dependence of the Hamiltonian in a rotating frame. Here both fields are assumed to be blueshifted with respect to resonance. The master equation for the reduced density matrix governing the internal and external dynamics reads ($\hbar=1$):

$$\begin{aligned}\dot{\hat\rho} &= -i\left(\hat H_{\rm eff}(\hat z)\hat\rho - \hat\rho \hat H^\dagger_{\rm eff}(\hat z)\right) + \mathcal{L}\hat\rho\\ \hat H_{\rm eff}(\hat z) &= \frac{\hat p^2}{2M} + V_{AL}(\hat z)\\ V_{AL}(\hat z) &= \sum_{k=1,2}\left(U_k - i\frac{\gamma_k}{2}\right)[a^k_-(\hat z)|-1\rangle + a^k_+(\hat z)|+1\rangle][(a^k_-(\hat z))^*\langle -1| + (a^k_+(\hat z))^*\langle +1|]\end{aligned} \tag{2}$$

In writing Eq. (2) we have adiabatically eliminated the excited states, which is valid in the low–saturation limit. $\hat H_{\rm eff}$ is an effective Hamiltonian, $\hat V_{AL}$ is the optical potential describing the interaction between the cooling laser and the atoms, M is the mass of the atom, and $|m\rangle$ are groundstate Zeeman–sublevels. The optical potential strengths U_i and the optical pumping rates γ_i are related to the detunings $\Delta_i=\omega^{\rm L}_i-\omega^{\rm A}_i$, the Rabi frequencies Ω_i and the atomic decay–rates Γ_i by $U_i=\Delta_i s_i/2$ and $\gamma_i=\Gamma_i s_i/2$, where $s_i=\frac{1}{2}\Omega_i^2/(\Delta_i^2+\Gamma_i^2/4)$ is the saturation parameter of the i–th transition. $\mathcal{L}\hat\rho$ is the standard recycling term.

As a first attempt to obtain some qualitative understanding of our system we calculate the adiabatic potentials. To this end we locally diagonalize the optical potential V_{AL} for $\gamma_i=0$. In a semiclassical picture this corresponds to the potentials seen by slow atoms. In the case of $J_g=1$ the two adiabatic potentials corresponding to the maximally and minimally coupled internal atomic states can be analytically calculated and are given by (see Fig. 1):

$$V_\pm(z) = (U_1+U_2) \pm \sqrt{U_1^2+U_2^2+2U_1U_2\cos\left(2z(k_1-k_2)\right)} \tag{3}$$

In the case of blue detuning of both laser frequencies relative to the respective transition–frequencies we find that the minima of the adiabatic potential $V_-(z)$ coincide with points which locally decouple from the interaction with the lasers (=zero lightshift). The period of the potentials is determined by the difference of the wavevectors $L_s=\pi/|\vec k_1-\vec k_2|$. The third adiabatic potential corresponding to the $m=0$ state shows no spatial modulation; it is, however, depleted by optical pumping in the initial stages of the cooling process and hence plays no important role (in 1D).

In a semiclassical picture (point particle with respect to the external degrees of freedom) atoms in the superposition state corresponding to the "darker" adiabatic potential would be drawn towards regions of minimal energy (dark areas); atoms in the orthogonal bright state would be repelled from these areas. As the magnitude of these forces is basically the same and the atom spends most of the time in the quasi–dark states, one gets a net trapping force towards the dark areas on average. In addition the spatial variation of the decay rate associated with the adiabatic potential $V_{+}(z)$ exhibits the proper trend needed for Sisyphus cooling, i.e. higher energies correspond to higher decay rates. Hence in addition to randomly jumping into the darker low lying states, the atoms loose energy by the spatially modulated hopping between the two optical potentials. Note that a combination of dark state cooling and Sisyphus cooling has been discussed by Marte *et al.* [7]. Such precooling mechanisms are important as the time scale of dark state cooling can be slow especially in two and three dimensional setups.

In order to get some quantitative answers about this trapping and cooling scheme, we determine the eigenvalues and eigenvectors of the full Hamiltonian including external momentum by numerical solution on a discrete spatial grid. In this case we can limit ourselves to one superperiod L_s. In a superlattice picture each state obtained above would correspond to a whole energy band. As the (tunnelling–) coupling between neighboring wells in our case is in general very small, we can treat them as isolated states. We have checked this numerically by extending the spatial grid to several superperiods. The corresponding energy shifts turn out to be very small.

Similar to the case of a dark FORT mentioned before, due to the finite size of the dark areas in space, we cannot expect to find eigenstates of the Hamiltonian, which are exactly decoupled from the light field (i.e. velocity selective dark states). Nevertheless we find, that the lowest lying states are (periodically) well localized in the dark regions and contain only very small contributions from the upper atomic states. In Fig. 1a we plot the energetically lowest eigenfunctions corresponding to each respective adiabatic potential. As expected, we find a strong localization of the lowest eigenfunctions in the dark areas, whereas the lowest eigenfunctions corresponding to $V_{+}(z)$ are bound to different areas of the superlattice. The spectrum of eigenvalues and the effective widths of the corresponding levels are shown in Fig. 1b. The parameters are as in Fig. 1a. The eigenvalue spectrum has two branches: a lower branch corresponding to weakly coupled states (labeled **dark**) and an upper branch of **bright** states. This notation becomes more obvious, if on looks at their corresponding effective widths as plotted using dashed lines. Relating these two branches to the corresponding eigenvectors we see, that as expected one gets a strong localisation of the low energy eigenstates towards the dark areas, while the bright states are bound to different spatial areas in the superlattice. Depending on the relative magnitude of the two laser fields the spatial extension of the lowest eigenstates can be less or much larger than an optical wavelength. Of course besides these localized bound states at energies above the well depth, we find 'free' delocalized solutions describing fast atoms. Note, that we have restricted our numerical

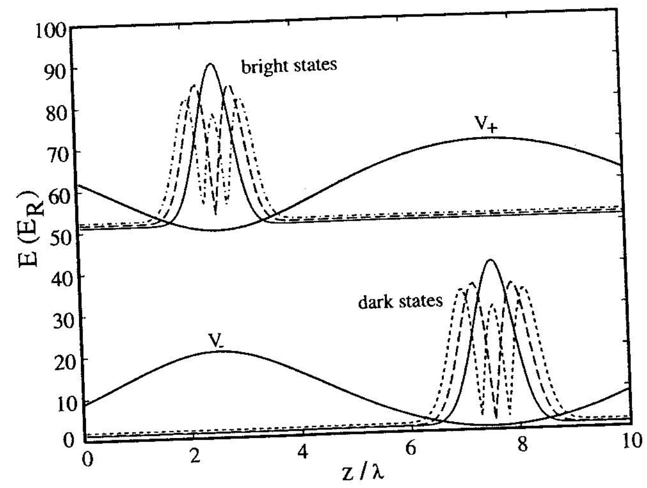

FIG. 1

Fig. 1a. Modulus of the three lowest eigenfunctions in position representation corresponding to the adiabatic potentials $V_{\pm}(z)$ (together with the adiabatic potentials) plotted over one superperiod L_s. The vertical position indicates the corresponding eigenenergy in units of the recoil energy $E_R = \hbar^2 k_1^2/2M$. The parameters are $k_2 = 1.05 k_1$, $U_1 = 25 E_R$ and $U_2 = 10 E_R$.

analysis to a single superperiod. As mentioned above in a larger spatial volume, each eigenstate would correspond to a band of nearly degenerate eigenstates. For most parameters the well to well coupling is, however, very small, so that each well can be treated separately.

3. Cooling and trapping dynamics

The decisive question regarding the usefulness of the proposed scheme is related to the distribution of the steady state population of the lowest localized quasi–dark levels, which determines the mean energy and the density. To calculate this we have to include optical pumping in addition to the coherent laser–field dynamics. This leads to a finite lifetime of the individual states and thus to a dynamical redistribution of the population among the various eigenstates. We use a similiar approach as in Refs. [8], and calculate the steady state population distribution using two methods: firstly a rate equation approach based on transition matrix elements between the different eigenvectors and secondly a standard Quantum Monte Carlo technique[9]. As the optical pumping rates γ_i depend on the spontaneous linewidth of the upper state as well as on the detuning and intensity of the laser fields and can thus be suitably tailored to a large extend. Fortunately, in the interesting parameter ranges both methods yield essentially the same

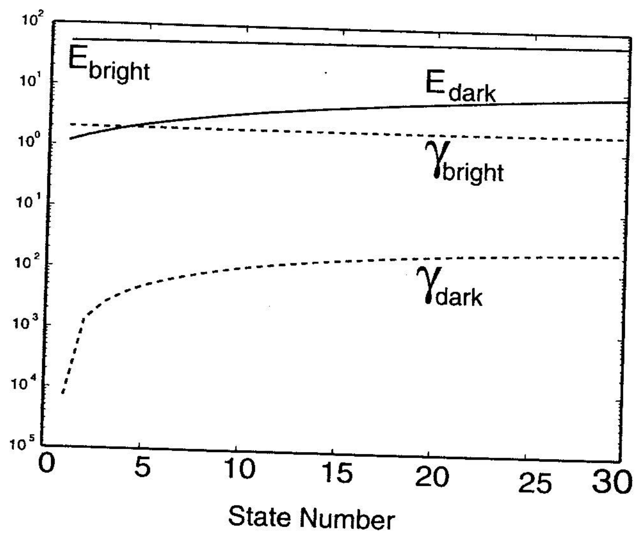

Fig 1b. Energy (solid line) and effective width (dashed line) of the atomic eigenfunctions. The vertical axis gives the eigenenergy in units of the recoil energy $E_R = \hbar^2 k_1^2/2M$ and the linewidth in units of the optical pumping rate . The parameters are $k_2 = 1.05k_1$, $U_1 = 25E_R$ and $U_2 = 5E_R$ and $\gamma_2 = \gamma_1/1000$.

promising result. Under suitable operating conditions most of the atomic population is concentrated in the lowest dark states and confined to a rather small volume in space. This is shown in Fig. 2, where we plot the steady state population of the ground state Π_0 for different various values of the optical pumping rate γ_2 as a function of the optical potential depth U_1. The other parameters are chosen as $U_2 = 3E_R$, $\gamma_1 = 1E_R$, $k_2 = 1.05k_1$, where $E_R = \hbar^2 k_1^2/2M$ is the recoil–energy.

Note that the best results are obtained for $U_1 > U_2$ and $\gamma_1 \gg \gamma_2$, which corresponds to the case of one nearly resonant laser and one strongly detuned and intense laser. The range of U_1 for which on obtains a large ground state population becomes broader for smaller values of the optical pumping rate γ_2. For $\gamma_2 = 0$ the population of the ground state increases monotonically with U_1. (Note that even in the case $U_1 \gg U_2$ we still have a modulation of the adiabatic potentials with a depth of the order U_2 and do not recover the standard velocity–selective coherent population trapping.) Although in this regime the highest ground state populations may be achieved, the system becomes very sensitive to γ_2. This behavior is caused by the fact, that for $U_1 \gg U_2$ the eigenfunctions are primarily determined by the part of the optical potential V_{AL} corresponding to the first pair of lasers. Therefore, the wavefunction of the (dark) ground state will be out of phase with respect to the second laser field. Even a small optical pumping rate γ_2 is sufficient to destroy the darkness of this state. By choosing the two optical potential

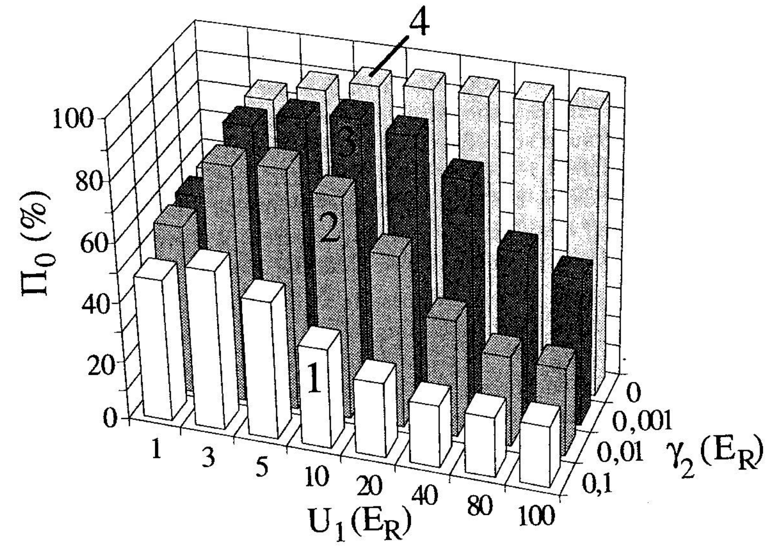

FIG. 2

Fig. 2. Steady state ground state population Π_0 as a function of U_1 for various values of γ_2. The other parameters: $U_2 = 3E_R$, $\gamma_1 = 1E_R$ and $k_2 = 1.05k_1$.

depths of the same order of magnitude, the ground state wavefunction can adjust to the two laser fields simultaneously.

Note that the probability to find the atom in the lowest eigenstate can reach almost unity for parameters which seem well in reach of present experimental capabilities. Furthermore we find, that in order to obtain a high Π_0, it is not important to have an extremely dark ground state, the relative magnitude of the lifetimes of the lowest states is essential. On the other hand in order to minimize atom–atom interactions in the multiatom case, a very dark lowest level should be advantageous. Such a very long lifetime of the ground state compared to all other states can be obtained in two limits: (a) the strong confinement limit, where the center of mass wavefunction is of the order of or narrower than an optical wavelength and (b) the case of a wide shallow potential, where the lowest eigenstate very much resembles a free (plane wave) dark state with a slow modulation of its spatial amplitude. Here one recovers standard VSCPT of free atoms. The parameters considered here correspond generally to the former case (a). For example, the width of the ground states for the parameters labeled by 1 – 4 in Fig.

2 is given by $\Gamma_1 = 2 * 10^{-3}$, $\Gamma_2 = 2 * 10^{-4}$, $\Gamma_3 = 7 * 10^{-5}$, and $\Gamma_4 = 5 * 10^{-5}$ in units of $E_R/\hbar$.

A further important quantity is the timescale of the cooling process. In the regime where the rate equations are valid, the cooling time scales linearly with the inverse optical pumping rates, provided the ratio γ_1/γ_2 is constant. A numerical analysis of the cooling time into the ground state based on the rate equations for the same parameters as above (see Fig. 2, labels 1 – 4) yields $\gamma_1\tau_1^{cool} \simeq 150$, $\gamma_1\tau_2^{cool} \simeq 500$, $\gamma_1\tau_3^{cool} \simeq 675$, and $\gamma_1\tau_4^{cool} \simeq 700$. We define the cooling time τ^{cool} as the time necessary to reach $1/e$ of the stationary ground state population, starting from a thermal distribution with a mean value of $20E_R$. These estimates show, that the cooling can be very effective.

Of course there exists a variety of possible atomic systems covering a wide range of parameters to be used in a corresponding experimental setup. An example for such systems are the $J_g = 1$ to $J_{e_1} = 0$, $J_{e_2} = 1$ transitions in metastable He, which exhibit an energy spacing of 30GHZ and hence a superlattice period of 0.5cm. In this case , and possibly using a MOT as an initial precooling mechanism, we could expect to confine a large fraction of the precooled atoms within a volume of only a few cubic wavelengths. To achieve this, effective light shifts U_i of the order of $10 - 1000E_R$ are sufficient. Of course in this case multi–atom effects can no longer be neglected, which will then give an upper limit for the achievable density. As the trapped states contain only a very small fraction of upper state population, this density can be expected to be very high.

A second example is ^{87}Rb, where one has two $F = 1$ to $F = 1$ transitions in the two finestructure components $5P_{1/2}$ and $5P_{3/2}$, separated by 15nm. Here the frequency splitting is much larger and the superperiod amounts to 26 wavelengths. The fact that this is not a closed transition will require a repumper as usual in alkali cooling experiments. An even shorter superperiod could be achieved using the 5P and 6P $F = 1$ sublevels, leading to superperiod of the order of a wavelength. In this limit the system is similar to a setup proposed by Grynberg and Courtois [10] using a magnetic field to create localized dark states. However, in such a limit starting from the initial density of a MOT, it seems hard to get more than one atom into a single superwell.

As a major drawback of our theoretical description, we have not included any many-atom effect in our dynamics. This would, however, seem a necessary requirement to obtain reliable answers about densities and temperatures in such a system. On the other hand, we have shown that the atoms are cooled down to relatively dark atomic states with most of the population trapped in the ground state of the system. So photon scattering and reabsorption, which is the main limiting factor for the density in a MOT[11] can be reduced by several orders of magnitude, while the confinement force can still be very strong. Similarly dipole-dipole interaction is very weak for the (quasi)dark states. As has been shown [12] significant effects only show up at mean atomic distances much less than the optical wavelength, corresponding to very high densities. If the proposed trap was loaded from a MOT, one could expect the global density (i.e. the density averaged over a superwell) to be of the order of or lower than the density in the MOT. At the minima of the adiabatic potentials, however, we expect a high *local* density of atoms. Of course effects related to collective radiative phenomena might have to be considered as well [13]. A study of such interatomic correlation effects might be a major possible application of such a dark and dense purely optical trap.

In this work we have restricted ourselves to a $J_g = 1$ ground state. An analogous scheme is also realizable in other atomic transition schemes which possess dark states. In particular we have done calculations for $J_g = 2$ and $J_g = 3/2$ ground states. Note that similar to standard VSCPT, our scheme can be generalized to 2D and 3D. For example, using a tetragonal configuration consisting of four pairs of laserbeams [14] one finds (dark) local minima of the adiabatic potentials. However, explicit numerical calculations for realistic parameters in 3D currently seem impossible. Finally we want to mention that a similar behavior also turns up by applying the two pairs of lasers to the same atomic upper state (bichromatic field) [15]. The quantitative results in this case are, however, a bit less promising.

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References

[1] W.D. Phillips: *Fundamental systems in Quantum Optics*, ed. by J. Dalibard, J.-M. Raimond and J. Zinn-Justin (North Holland, Amsterdam, 1992). For an overview of laser cooling see, e. g., C. Cohen-Tannoudji: *ibid.*;

[2] M. Kasevic: *private Communication*;

[3] W. Ketterle, D. B. Kendall, J. A. Michael: *Phys. Rev. Lett.* **70**, (1994) 2253;

[4] A. Aspect et al.: *Phys. Rev. Lett.* **61** (1988) 826;

[5] J. Lawall et al.: *Phys. Rev. Lett.* **73** (1994) 1915; M.A. Ol'shanii, V. Minogin: *Quantum Opt.* **3** (1991) 317;

[6] J.D. Miller et al.: s Phys. Rev. A **47** (1993) R4567; R. Dum, P. Marte, T. Pellizzari, P. Zoller: *Phys. Rev. Lett.* **73** (1994) 2829;

[7] P. Marte, R. Dum, R. Taieb, P. Zoller: *Phys. Rev. A* **49** (1994) 4826;

[8] Y. Castin, J. Dalibard: *Europhys. Lett.* **14** (1991) 761; P. Marte, R. Dum, R. Taieb, P. Zoller: *Phys. Rev. A* **47** (1993) 1378;

[9] see e.g. J. Dalibard, Y. Castin, K. Molmer: *Phys. Rev. Lett.* **68** (1992) 580; H. Carmichael: *LNIP* (Springer, New York, 1994), Vol. m18; R. Dum, P. Zoller, H. Ritsch: *Phys. Rev. A* **45** (1992) 4879;

[10] G. Grynberg, J.-Y. Courtois: *Europhys. Lett.* **27** (1994) 42;

[11] K. Ellinger, J. Cooper, P. Zoller: *Phys. Rev. A* **49** (1994) 3909;

[12] E. Goldstein et al.: *Appl. Phys. B*, in press;

[13] see dicussion in section V of: I. Cirac, M. Lewenstein, P. Zoller: *Phys. Rev. A* **50** (1994) 3409;

[14] P. Verkerk et al.: *Europhys. Lett.* **26** (1994) 3; *ibid* 171;

[15] Yu. B. Ovchinnikov, R. Grimm, A. I. Sidorov: *Opt. and Spectr.* **76** (1994) 88;