Synchronization of atomic quantum systems in multi-site optical trapping potentials

Malte Schlosser, Jens Kruse, and Gerhard Birkl

Institut für Angewandte Physik, Technische Universität Darmstadt,
Schloßgartenstraße 7, 64289 Darmstadt, Germany

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Abstract

Advanced quantum technologies, such as quantum simulation, computation, and metrology are thriving for the development of large-scale configurations of identical quantum systems. Atomic and molecular systems have the advantage of identical intrinsic properties but need to be placed in identical environments as well. In this letter we present an experimental scheme for compensating the differential Stark shift induced by optical dipole traps. The compensation results in a strong suppression of dephasing effects and leads to an increase of the observable coherence time. This behavior is achieved by the addition of a second, near resonant light field to the trapping field. The achieved suppression of dephasing is demonstrated by analyzing the hyperfine state phase evolution of atomic ensembles of $^{85}$Rb trapped in a two-dimensional array of dipole traps via Ramsey spectroscopy. The experimental method presented here is expandable to all atomic species trapped in various dipole trap configurations of arbitrary wavelength.

*gerhard.birkl@physik.tu-darmstadt.de
INTRODUCTION

In recent years there has been an enormous progress in the field of quantum engineering with ultracold neutral atoms trapped and manipulated by optical means. Applications in quantum metrology and quantum information science, such as all-optical atomic clocks or quantum information processing, benefit from the extraordinary extend of control provided by light fields with respect to the internal and external degrees of freedom of neutral atoms. One main requirement in most of these experiments is the capability to precisely manipulate the electronic states of an atomic system which features long coherence times as well [1]. Unfortunately, in most quantum systems the influence of optical potentials is state dependent and thus perturbs a coherent evolution of quantum phases due to differential effects [2–5]. Several approaches to separate internal and external dynamics have been made. Specially the realization of trapping atoms in potentials of a so called ”magic wavelength” represents a break through in high precision spectroscopy like for example in optical clocks [6]. A magic wavelength appears in some elements whose complex electronic structure results in the cancellation of differential effects for certain wavelengths of the trapping fields, thus yielding a vanishing perturbation of the atomic phase evolution. Although there are no nondisturbing optical potentials for alkali atoms generated by light of simple magic wavelength [7], these elements still play a crucial role in quantum optics, representing a well determined bosonic test bed. One option to minimize the spatial dependence of the light field for alkalis is for example the trapping of an atomic ensemble in a far blue detuned dipole trap, where the atoms spend most of their time in a nondisturbed space [8]. Other possibilities are to use a vector potential of partly circular polarized light to compensate for differential light shifts [9] or to suppress dephasing mechanisms by a sequence of microwave pulses [10].

Here we present the experimental realization of a simple and universally extendable scheme of compensating the differential light shift induced by dipole traps. We perform Ramsey spectroscopy on $^{85}$Rb trapped in an array of dipole potentials to determine the coherence time. By using a second, near resonant light field we can compensate the differential light shift (Fig. ??) and strongly suppress atomic dephasing [11]. Due to this we increase the coherence time which is one of the most important key elements in quantum metrology and quantum information systems.
universal method for the compensation of the differential ac-stark effect

The interaction of a laser field with frequency $\omega_L$ and intensity $I$ and an atom with internal states $|g_i\rangle$ and $|e_j\rangle$ can be described by the dipole potential. In the rotating-wave-approximation (RWA) the ground state ac-stark shift a multi-level atom experiences is given by

$$\Delta E_i(r) = \frac{3\pi c^2 \Gamma^2}{2 \omega^3} I(r) \sum \frac{c_{ij}}{\Delta_{ij}} .$$

(1)

Here $\Gamma$ is the natural linewidth and $\omega$ is the atomic transition frequency. The summation takes into account the contributions for the different coupled excited levels $|e_j\rangle$, each with its respective transition coefficient $c_{ij}$ and detuning $\Delta_{ij} = \omega_L - \omega_{ij}$. Depending on the sign of $\Delta$ the dipole potential is either repulsive ($\Delta$ positive) or attractive ($\Delta$ negative). Most atoms like the alkali elements have more than one long living ground state due to their nuclear spin. These hyperfine states, defined as $|\|0\rangle$ and $|\|1\rangle$, are used in several applications in metrology like frequency standards in atomic clocks or as qubit basis in quantum computing. One major drawback of the widely used alkali elements is the nonnegligible influence of the dipole trap on the hyperfine splitting of the ground state. Although the dipole matrix elements for the two ground states are equal for linearly polarized light, the difference in detuning caused by their frequency splitting $\omega_{HFS}$ yields a differential light shift

$$\hbar \delta_{\text{diff}}(r) = \frac{3\pi c^2 \Gamma^2}{2 \omega^3} \left( \sum \frac{c_{0j}}{\delta_{0j}} - \sum \frac{c_{1j}}{\delta_{1j}} \right) \cdot I_{\text{trap}}(r) .$$

(2)

This light shift $\delta_{\text{diff}}(r)$ leads to an intensity dependent shift of the hyperfine splitting. A variation in the trap depth modifies the observed frequency shift. Regarding the energy distribution of a thermal atomic ensemble in a dipole trap, described by the three-dimensional Boltzmann distribution with probability density

$$p(E) = \frac{E^2}{2(k_B T)^3} \exp \left( -\frac{E}{k_B T} \right) ,$$

(3)

the differential light shift also results in a broadening of the averaged atomic transition. The probability density $p(E)$ can be rewritten as $p(\delta_{\text{diff}})$ converting the energy distribution into a frequency spread. This frequency broadening induces a dephasing of the atomic state evolution and limits the spectroscopic resolution. In summary the differential light shift causes two major problems. First the effective resonance frequency of the clock transition
varies from trap to trap depending on the dipole trap depth. Secondary the measurement 
time and thus the frequency resolution is limited by dephasing mechanisms due to the 
differential light shift.
To compensate for the differential light shift induced by the dipole potential a second light 
field, coupling the two hyperfine ground states to one of the excited fine structure states, is 
introduced. The compensation field is detuned in between the two hyperfine ground states 
with $\Delta_c \approx -\omega_{\text{HFS}}/2$ compared to the state $|0\rangle$ and $\Delta_c \approx +\omega_{\text{HFS}}/2$ to the state $|1\rangle$ yielding 
in a differential shift of

$$
\hbar\delta_c(r) = \frac{3\pi c^2 \gamma}{2\omega^3} \left( \sum_{2c_{0j}} \frac{2c_{0j}}{\omega_{\text{HFS},j}} + \sum_{2c_{1j}} \frac{2c_{1j}}{\omega_{\text{HFS},j}} \right) I_c(r).
$$

(4)

The index $j$ in the detuning accounts for slight modifications caused by the hyperfine man-
ifold of the excited state. Due to the change of sign in the detuning the additional laser 
field increases the effective hyperfine splitting. This behavior is contrary to the compression 
of the atomic resonance by a far red or blue detuned dipole trap. By simply adding both 
light fields one can compensate for the differential light shift. For this compensation only 
a balancing of the differential effect is necessary, which is orders of magnitude smaller than 
the ac stark shift $[12]$. As a result the modification of the absolute trap depth is negligible. 
Nevertheless the ac stark shift of both hyperfine states becomes equal. To cancel the inho-
mogeneous broadening all over the dipole potential simultaneously the two light fields have 
to be perfectly mode matched

$$
\delta_{\text{diff}}(r) + \delta_c(r) \overset{!}{=} 0 \quad \rightarrow \quad I_c(r) = \eta \times I_{\text{trap}}(r)
$$

(5)

where the ratio $\eta$ approximately scales with $\eta \approx (\omega_{\text{HFS}}/(2\Delta)) k^2 [11, 13]$. 
Although the detuning $\Delta_c$ is in the range of the ground state hyperfine splitting the scattering 
rate of both laser fields ($\Gamma_c$ and $\Gamma_{\text{trap}}$) is in the same order of magnitude $\eta_{sc} = \Gamma_c/\Gamma_{\text{trap}} \approx 1$, 
due to the small amount of required compensation power. We point out, that for optimal 
compensation this ratio is always valid and does not depend on the detuning of the dipole 
laser or the alkali element.
EXPERIMENTAL DEMONSTRATION OF MAGIC-WAVELENGTH BEHAVIOR FOR A SCALABLE 2D ARRAY OF OPTICAL DIPOLE TRAPS

In the experiment described here we work with ultracold $^{85}\text{Rb}$ atoms trapped in an array of well separated dipole traps. Our system represents a scalable architecture for a quantum register using the two hyperfine levels of $^{85}\text{Rb}$ $|F = 2, m_F = 0\rangle = |0\rangle$ and $|F = 3, m_F = 0\rangle = |1\rangle$ as the qubit states and considering a single atom in each trap as a qubit [14]. We perform single qubit rotations by coupling the two clock states with a coherent two photon transition.

The dipole trap array is created by illuminating an array of microlenses with a Gaussian laser beam [15]. By imaging the foci of the focal plane into our vacuum chamber we superimpose the dipole trap array with the cold atomic cloud of $^{85}\text{Rb}$ cooled in a MOT. In the experiments presented here we load each of the traps with several tens to hundreds of atoms during a sequence of optical molasses. All traps have a waist of $w_{\text{trap}} = 3.7 \pm 0.1 \, \mu\text{m}$ ($1/e^2$-radius) and the trap separation is 55 $\mu\text{m}$. As a dipole trap laser we use a Titanium Saphire laser (TiSa) at a wavelength of $\lambda_{\text{TiSa}} = 810.1 \, \text{nm}$. The laser is intensity stabilized ($\Delta P/P \approx 10^{-4}$) at a power of $P_{\text{TiSa}} = 41 \pm 1 \, \text{mW}$ with a waist of $w_0 = 500 \pm 3 \, \mu\text{m}$ at the position of the microlens array. This results in a power of $P_{\text{ct}} = 0.72 \, \text{mW}$ in the central trap leading to a trap depth of $U_{\text{ct}} = k_B \cdot 40 \, \mu\text{K}$. The trapping frequencies are $\Omega_r = 2\pi \cdot 5.5 \, \text{kHz}$ in radial direction and $\Omega_a = 2\pi \cdot 270 \, \text{Hz}$ in axial direction. We have the ability to resolve all single traps in parallel with a CCD-camera by fluorescence imaging and we showed elsewhere that we can address single traps with a tight focused laser beam [14] or a spatial light modulator [16].

Precision spectroscopy by performing Ramsey spectroscopy and its extension the spin echo technique determines the clock transition frequency and the coherence time of the system. The coupling of the two hyperfine states is realized by a pair of phase locked lasers with detuning $\delta_{RL}$. From the measured oscillation frequency of the Ramsey signal we determine a shift of the effective atomic resonance of trapped atoms compared to the unperturbed hyperfine splitting $\omega_{\text{HFS}}$: $\delta_{\text{shift}} = \omega_{\text{HFS}} - \delta_{RL}$. One contribution to the observed shift is the difference of the second order magnetic field offset $\delta_z = 264 \pm 6 \, \text{Hz}$ of the two clock states. This offset occurs due to the applied magnetic bias field of $B = 45.2 \pm 0.5 \, \mu\text{T}$, which is assumed to be equal for all traps. The second main contribution of the frequency shift
is the differential light shift induced by the dipole trap. By using Ramsey spectroscopy the differential light shift of all traps can be measured in parallel. In Fig. ?? (triangles) the measured transition frequency of 54 separated atomic ensembles trapped in the dipole trap array is presented. The transition frequencies are shown relative to the unperturbed resonance \( \omega_{HFS} \), corrected by the magnetic shift \( \delta Z \). The observed differential light shift reflects the spatial intensity profile of the Gaussian dipole beam at the microlens array. The decreased transition frequency of deep center traps compared to the small shift of more shallow traps leads to a frequency spread of 260 Hz regarding all trapped atomic ensembles. For the compensation of the differential light shift a coupling to the D1- or D2-line of the additional light field is feasible to realize the proposed scheme. However, to minimize variations in the spatial overlap due to chromatic aberrations we choose the D1-line. To implement the scheme we use a second TiSa laser, frequency stabilized at a wavelength of \( \lambda_c = 794.978 \) nm, thus centered in between the ground state hyperfine levels. Both laser fields are spatially mode matched by sending them through one polarization maintaining single mode optical fiber. We optimized the laser power \( P_c \) required for optimal compensation by experimentally minimizing the frequency spread between different dipole traps. In Fig. ?? (circles) the light shift of 54 analyzed atomic ensembles in the dipole trap array with the additional compensation field is presented. Compared to the differential light shift induced by the single dipole beam, the frequency spread between the compensated traps is reduced by a factor of 20 with an average value close to the non shifted resonance. We stress the point that both experiments have been carried out under identical parameters, except the application of the compensation laser with a measured power of \( P_c = 4.3 \pm 0.2 \) nW in front of the microlens array. As seen in Fig. ?? the differential light shift is slightly overcompensated and an opposite frequency shift is observable. To analyze the limitation of the compensation of the differential light shift Fig. ?? shows the residual frequency spread after optimizing the compensation power to \( P_c = 4.2 \pm 0.2 \) nW. This results in a ratio between both laser fields of \( \eta_{\text{meas}} \approx 1.02 \pm 0.6 \times 10^{-7} \).

The residual structure remaining in the compensated frequency spread arises due to a slight displacement between the two laser fields of 810 nm and 795 nm at the position of the microlens array. By simulating a misalignment between both beams of 8 \( \mu \)m the residual frequency structure is fitted well. The displacement of the two beams can be explained by a slight decentered (200 \( \mu \)m) collimation lens behind the optical fiber. Analyzing the differ-
ence between the two curves the resulting residual frequency spread decreases to $\sigma = 1.7$ Hz between all analyzed dipole traps.

The compensation field not only equalizes the observed hyperfine splitting in different dipole traps, it also effects the dephasing of each single atomic ensemble. In Fig. ?? the Ramsey oscillations of a single trap without (top) and with (bottom) the application of the compensation laser are shown. In the latter case the decay of the Ramsey amplitude is strongly suppressed and the inhomogeneous dephasing time $T^*_2$ increases by a factor of 13. After a measurement time of 120 ms there is still fringe contrast observable.

To determine the coherence time limited by irreversible dephasing we expand the Ramsey measurement to an echo experiment. Due to this technique any symmetric dephasing caused by static inhomogeneities is reversed by the inverting mirror $\pi$ pulse and thus not leading to a decay in the oscillation amplitude. In our setup irreversible dephasing is mainly caused by heating, leading to a change in the vibrational state of the atoms in the trap during the measurement time $t$. Due to the differential light shift this energy change results in a non-reversible broadening of the atomic resonance. The decay of the spin-echo amplitude can be described by the irreversible dephasing time $T'_2$, defined as the $1/e$ decay time $[17]$. On the basis of a life time of $t_l = 1.2 \pm 0.1$ s the calculated heating rate leads to an theoretical dephasing time of $T'_2 = 50 \pm 10$ ms. In the uncorrected dipole trap array the measured $T'_2$-time evaluates to 60 ms. The application of the compensation laser results in a vanishing sensitivity to dephasing mechanisms caused by changes in the vibrational state, verified by the extended $T'_2$-time of 87 ms. Our technique of realizing dipole traps without differential effects by simple means yields an 45%-increase of the coherence time. In the case of an improved compensation with optical components optimized for the two different wavelength an extension of the coherence time up to the limit of spontaneous scattering processes should be possible.

**CONCLUSIONS**

In this article we presented an universal scheme of compensating the differential light shift utilizing a second light field detuned in between the two ground state hyperfine levels. Due to the compensation we equalized the transition frequency in an array of dipole traps varying in depth. Because of the spatial overlap between both light fields dephasing mechanism induced
by spatial inhomogeneities of the dipole potential are strongly reduced and an increase of the coherence time of 45% is observed. The presented compensation of differential effects is limited by chromatic aberration, which can be reduced by optimized optical alignment. The demonstrated simple scheme is expandable to all alkali atoms and beyond and induces the same amount of spontaneous scattering as the trapping laser, independent of its actual detuning.

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