Giant light shift of atoms near optical microstructures

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Atoms coupled to optical fields confined in one and two spatial dimensions in solid state microstructures can experience very large light shifts if the driving frequencies are close to a resonance of the microstructures and an atomic transition. Using the simple example of a quasi one-dimensional waveguide structure we can analytically calculate the atomic AC Stark shift and the modifications of the light field induced by the presence of the atom. A large enhancement of the effective interaction strength is found due to a non uniform mode density. Experimentally this should be visible by monitoring the scattered light field as well as by the modification of the atomic trajectories bouncing from the evanescent light.

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In the past years we have seen spectacular advances in our ability to cool atoms to nanokelvin temperatures and control their motional degrees of freedom down to the quantum level \[.\] In parallel, the miniaturization of optical microstructures has reached the level where fabrication almost at the atomic scale is feasible \[.\] One of the great goals in the near future is to bring together these technologies in a generation of integrated optical quantum devices \[.\] As a central point to utilize such devices, we must understand the behaviour and transport of the atomic matter subject to subwavelength-structured electromagnetic radiation fields. Evanscent fields created by tailored dielectric microstructures are of primary interest in this research. A series of experiments have demonstrated the usefulness of evanescent optical fields to realize atom mirrors \[] or quasi 2D surface traps \[.\] However, these surface setups have all been based on macroscopic fields regarding the transverse spatial dimension and photon numbers involved. Hence a single particle has almost no effect on the field and a single photon field would give only a negligible force on an atom. In this Letter we show that the effective atom-photon interaction can strongly be enhanced for atoms at the surface of dielectric microstructures, where the field is partly confined in some directions. Note that for a full 3D confinement one would recover a setup as in cavity QED \[] with special boundaries. Here we reveal new surprising phenomena beyond the effects of a single strongly coupled radiation mode. For example, a continuum of travelling wave modes can induce a huge atomic AC Stark shift exceeding by orders of magnitude the natural linewidth. In conjunction with the large light shift, a strongly increased photon scattering rate by a single atom takes place, which could be used for position and state selective single atom detection and manipulation schemes \[.\]

Atomic light shift in 1D continua of modes.- Let us consider a two-level atom with resonance frequency \(\omega_a\) and free-space spontaneous emission rate \(\Gamma\) in (or close to) a dielectric medium with refractive index \(n_0\) which is assumed to be infinitely extended into the \(z\) direction, but with a transverse dimension of the order of an optical wavelength. This microstructure supports optical modes which are described by annihilation operators \(a_n(k)\), where \(n\) labels the transverse mode index and \(k\) the longitudinal wave number. The annihilation and creation operators fulfill the standard commutation relation

\[
\left[a_n(k), a_{n'}(k')^\dagger\right] = \delta_{n,n'} \delta(k-k').
\]

The corresponding frequencies are denoted by \(\omega_n(k)\) and the mode functions read

\[
f_n(k,\mathbf{x}) = \exp(ikz)f_n^{(T)}(k, x, y).
\]

The mode functions are normalized such that

\[
A = n_0^2 \int_{A_1} dx dy |f_n^{(T)}(k, x, y)|^2 + \int_{A_2} dx dy |f_n^{(T)}(k, x, y)|^2,
\]

where the first integral goes over the part of the mode function inside and the second integral over the part outside the dielectric medium. \(A\) is the cross section of the microstructure. The positive frequency part of the electric field is then given by

\[
E^{(+)}(\mathbf{x}) = \sum_n \int dk E_0(\omega_n(k))f_n(k, \mathbf{x})a_n(k)
\]

with \(E_0(\omega) = \sqrt{\hbar \omega/(2e_0A)}\) the electric field of a single photon.

In the following we will assume that the medium is pumped by monochromatic light of frequency \(\omega_p \approx \omega_a\) and that only modes with frequencies close to this contribute to the system dynamics. We will thus replace \(E_0(\omega_n(k))\) by the corresponding value at the pump frequency \(E_0 = E_0(\omega_p)\) and pull it out of the integral in Eq. \[\]

In dipole and rotating wave approximation and in a frame rotating with \(\omega_p\) the total system dynamics is then governed by the Hamiltonian

\[
H = -\Delta_0 \sigma^+ \sigma - \sum_n \int dk \Delta_n(k)a_n^\dagger(k)a_n(k)
\]

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\]
Here $\sigma$ is the atomic lowering operator, $\Delta_a = \omega_a - \omega$, $\Delta_n(k) = \omega_p - \omega_n(k)$, $x_0$ is the position of the atom, and $g = \mu E_0$ where $\mu$ is the atomic dipole moment. The first line of Eq. (6) describes the free field evolution, the second line the atom-light coupling, and the third line the coherent pumping of the light modes.

Allowing for photon losses from the dielectric medium at a rate $2\kappa$ and assuming that the dynamics of the electric field occurs on a much faster time scale than the atomic dynamics, we may adiabatically eliminate the photon operators from the Hamiltonian. In the limit of small atomic saturation and neglecting quantum noise terms, we finally obtain the Heisenberg equation of motion for the atomic internal variable $\sigma$,

$$\frac{d}{dt} \sigma = (i\Delta_a - \Gamma + L)\sigma - \chi$$

where

$$L = g^2 \sum_n \int dk \frac{|f_n(k,x_0)|^2}{i\Delta_n(k) - \kappa}$$

$$\chi = g \sum_n \int dk \frac{f_n(k,x_0)\eta_n(k)}{i\Delta_n(k) - \kappa}.$$  

We therefore see that the coupling of the atom to the field modes gives rise to a driving term $\chi$ of the atomic operator $\sigma$ and, additionally, to an atomic light shift $L$. In general, $L$ is complex and the atom will thus experience a frequency shift according to the imaginary part of $L$ and a resonance line broadening described by its real part. Depending on the functional dependence of $\Delta_n(k)$, the light shift obtained from the coupling of the atom to continuous one-dimensional sets of electric modes can significantly alter the atomic dynamics, especially if $\Delta_n(k)$ is a highly asymmetric function. Since dielectric microstructures can be fabricated with high accuracy, this offers the possibility to tailor the atom-light interaction to a large extent.

Specific example of a microstructure. - In order to discuss this effect quantitatively in more detail, we will in the following concentrate on a specific example of such a microstructure.

Let us consider a dielectric medium with a rectangular cross section of height $D_x$ and width $D_y$ as sketched in Fig. 1, but infinitely extended in the $z$ direction. The height $D_x$ is chosen such that for large $D_y$ two transverse modes are supported at the pump frequency $\omega_p$ as indicated in the figure. (For simplicity we neglect polarization issues here and therefore assume a scalar electric field.) The dielectric surfaces perpendicular to the $y$ axis, on the other hand, are supposed to be coated with a highly reflecting metallic layer, and the width $D_y$ is chosen small enough such that only a single mode is supported in that direction. For given $D_y$ the wave number in $y$ direction is thus fixed to $k_y = \pi/D_y$ for all modes. Hence, for a given optical frequency, decreasing $D_y$ will reduce the wave number $k$ in $z$ direction. At a certain threshold width, $k$ will vanish and the mode is no longer supported for widths smaller than that threshold.

In Fig. 2 we show the numerically calculated frequencies and longitudinal wave numbers of the modes supported by such a structure. The width $D_y$ is chosen such that the threshold of the first excited branch coincides with the pumping frequency $\omega_p$, i.e., $\Delta_1(0) = 0$. Hence, for all of the modes in this branch $\Delta_1(k) \geq 0$ and a large atomic lightshift $L$, Eq. (6), can be expected. Moreover, at threshold we have $\frac{d\Delta_1(k)}{dk} = 0$ and therefore a large number of modes contributes nearly resonantly to the integral in $L$. Thus, the presence of the microstructure largely enhances the mode density near resonance.
1.05

\( L \) (units of \( \omega_p \))

\( \omega_{th} \) (units of \( \omega_p \))

FIG. 3: Light shift \( L \) versus threshold frequency \( \omega_{th} \). The solid line is \( \text{Re}(L_1) \), dashed line is \( \text{Re}(L_0) \), dash-dotted line is \( \text{Im}(L_1) \). The atomic parameters correspond to the \( D_2 \) line of Rb, \( \kappa = 0.001\omega_p \), \( n_0 = 1.5 \), \( D_s = \lambda_p/\sqrt{n_0^2 - 1} \). The atomic position is \( (x, y, z) = (D_x/2, D_y/2, 0) \), i.e., the point of maximum coupling at the surface of the dielectric medium.

An analytic approximation of Eq. (3) can be obtained assuming that the transverse part of the mode functions is given by a unique \( f_n^{(T)}(x, y) \) within each branch of modes for all the relevant frequencies, that is, assuming constant values of \( k_{xn} \) and \( k_{yn} \) and therefore a constant transverse wave number

\[
q_n = \frac{1}{n_0} \sqrt{k_{xn}^2 + k_{yn}^2}.
\]

Introducing an exponential convergence factor to cut off high frequencies which in the dipole approximation lead to an unphysical logarithmic divergence, we find by complex contour integration

\[
L_n = -2\pi \frac{q^2 n_0}{c} \left| f_n^{(T)}(x_n, y_n) \right|^2 \sqrt{1 - \left( \frac{q_n}{r_n + i s_n} \right)^2}.
\]

Here

\[
r_n = \frac{1}{c} \left[ -u_n + \sqrt{u_n^2 + \kappa^2 \omega_p^2} \right]^{1/2},
\]

\[
s_n = \frac{\kappa \omega_p}{c} \frac{1}{r_n},
\]

\[
u_n = 1 \left( c^2 q_n^2 + \kappa^2 - \omega_p^2 \right).
\]

For the parameters used in this Letter we compared these results with numerical integrations of Eq. (3) and found excellent agreement.

An example for the light shift \( L \) is depicted in Fig. 3. We plot the real part of the contribution \( L_0 \) from the lower energy branch of electromagnetic modes (the imaginary part being approximately zero) as well as the real and imaginary parts of \( L_1 \), the contribution from the excited states, versus the threshold frequency \( \omega_{th} \) of the excited branch. For the chosen parameters, \( \omega_{th} \) is a nearly linear function of the width \( D_y \), which varies approximately for 10% within the plotted range.

We note that \( L_1 \) is approximately constant and below one atomic linewidth. Hence the light shift induced by a branch of modes far above threshold is in fact insignificant. The second (near-resonant) branch of modes, on the other hand, yields a large light shift. For \( \omega_{th} < \omega_p \), travelling wave solutions exist in the excited branch at the pump frequency, and the light shift is dominated by its real part leading to increased spontaneous atomic decay by enhanced emission of photons into the 1D microstructure. For \( \omega_{th} > \omega_p \), no travelling solutions exist, \( L \) is imaginary, and the main effect is a shift of the atomic frequency. At threshold and assuming \( \kappa \ll \omega_p \), Eq. (4) can be approximated by

\[
L_n = 2\pi \frac{q^2 n_0}{c} \left| f_n(x_n) \right|^2 \frac{i - 1}{2} \sqrt{\omega_p/\kappa}.
\]

Hence, the maximum possible light shift is determined by the ratio of the photon loss rate from the microstructure to the optical frequency. Thus, for our specific example of a structure the limiting factor will be the reflectivity of the metallic coatings on two sides of the dielectric surface. For example, a reflectivity of 99% yields a decay rate \( \kappa \approx \omega_p/1000 \) as used in Fig. 3. Both the magnitude and the scaling of the calculated light shift is very different from the one observed between metallic plates [8].

Let us now address the question how this large atomic light shift could be observed experimentally. Since the whole effect is due to the strong coupling of the atom to the confined light modes, an obvious possibility is to
detect the backaction on the light field. Assuming that only a single travelling wave with wave number \( k_0 \) of the lower branch of modes is pumped and with the same simplifications as used to obtain Eq. (2) yields the following stationary electric field \( E(x) \):

\[
E(x) \propto e^{i k_0 z} f^{(T)}_0(x, y) - L_0 \frac{e^{i \omega_0 (ir_0 - s_0)}}{\Delta_0 - \Gamma + L} f^{(T)}_0(x, y) - L_1 \frac{e^{i \omega_0 (ir_1 - s_1)}}{\Delta_1 - \Gamma + L} f^{(T)}_1(x, y, y_a),
\]

where the first line gives the field of the single pumped mode, the second line is the field of the light scattered into the lower branch of modes, and the third line is the field scattered into the upper branch. As an example we plot the field intensity along the \( z \) direction at the center of the medium, Fig. 5(a), and on the surface, Fig. 5(b).

At the center of the structure all modes of the excited branch vanish and the electric field is formed by the lower branch only. Since the light shift \( L_0 \) according to this branch is small, the change of the electric field is small too. However, we see that the atom (at position \( z = 0 \)) scatters some light from the pumped mode into its degenerate counter-propagating mode. Hence, on top of the constant intensity of the pumped mode, there appears a standing wave structure on one side of the atom. Due to the damping of the light modes, this standing wave has an exponentially decaying envelope with a decay distance of \( d_0 = c/(n_0 K) \). The electric field at the surface, on the other hand, is dominated by the large light shift \( L_1 \) due to the excited branch of modes and therefore has a much larger change of amplitude, see Fig. 5(b). Similar enhancement of light scattering has been predicted for a dielectric wire in a metallic wave guide [6]. According to Eq. (14), we find again an exponential decay with an approximate decay length of \( d_1 = c/(n_0 \sqrt{\omega_0 p}) \) which is much shorter than that of Fig. 5(a). A spatially resolved detection of the photons lost through the coatings of the dielectric structure would thus reveal the significant change of the electric field intensity and would serve as an implicit measurement of the enhanced atomic light shift.

An alternative method to detect the light shift could come from an experiment where a cloud of cold atoms is dropped onto the microstructure and reflected by the evanescent light field. If the atomic cloud is dilute enough such that the mean distance between the atoms is larger than \( d_1 \), each atom will be scattered individually and the reflection of the cloud will essentially be specular. On the other hand if atoms are closer than \( d_1 \), they will interact with a distorted light field as shown in Fig. 5(b). Since the modulation of the light intensity along \( x \) is roughly of the order of the total intensity and since the periodicity is of the order of an optical wavelength, the forces along \( x \) will be comparable with the force in \( z \) direction. Hence the average reflection of a cloud of atoms will be highly diffusive in this regime.

**Conclusions.-** We have chosen here to discuss the interaction of an atom with a quasi 1D optical waveguide. However, our theory is easily applicable to atom-photon interactions in a broad range of micro-optical structures. Corresponding experiments can have impact on our basic knowledge about atomic structure and quantum electrodynamics [7]. On the practical side, the strongly enhanced atom-photon interaction could be the basis of a new generation of micro-optical devices involving only very few atoms and photons.

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