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A strong magneto-optical force exerted on neutral atoms

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Abstract. — We report the theoretical prediction and the observation of a stimulated novel magneto-optical radiation force, acting on atoms located in a static magnetic field and irradiated by the resonant field of two counterpropagating monochromatic laser beams with different directions of polarization.

Introduction.

In the presence of a static magnetic field the radiative force acting on a multilevel atom in a laser field may acquire new properties [1, 2]. Of particular interest for the creation of new types of atomic traps is a confining potential-like character of the radiation force. To produce a magneto-optical force (MO) of pseudopotential character [1], a scheme based on the spontaneous force has been so far used. In this case, the spatial dependence of this force is determined by the Zeeman-shift in the magnetic sublevels of the atom induced by a spatially-nonuniform static magnetic field. One of the short-comings of this scheme is that the maximum of the MO force is restricted by the upper limit of spontaneous force $\hbar k \gamma$, where $2 \gamma$ is the natural width of excited state.

Here we want to show that there is a MO force based on induced photon reemission between two counter-propagating light waves having different directions of polarization vectors.

It is essential that the value of the new MO force may considerably exceed $\hbar k \gamma$.

In section 1 this MO force is interpreted in terms of dipole force rectification [3]. The MO force acting on an atom with the transition $J = 0 \rightarrow J' = 1$ is calculated exactly here.

Section 2 gives a classical interpretation for the MO force and the underlying process of stimulated photon redistribution between the two traveling waves in terms of the Hanle effect.

Section 3 presents the results of an experiment where we studied the deflection of a sodium atomic beam by the MO force.

1. Magneto-optical force for an atom with $J = 0 \rightarrow J' = 1$ transition.

Our model system consists of atoms with a closed optical transition $J = 0 \rightarrow J' = 1$ in the field of two linearly polarized traveling waves with equal frequency $\omega$ and intensity $I_0$ counterpropagating along the z-axis (Fig. 1) [4]. We assume that the polarization directions of
these two fields are arranged at an angle of \( \varphi \) with respect to each other. A static magnetic field \( \mathbf{B} \) is applied in \( z \)-direction to split the Zeeman sublevels of the excited state \( J' = 1 \) (Fig. 2).

![Diagram of light and magnetic fields](image)

**Fig. 2.** — Our model system — an atom with a \( J = 0 \rightarrow J' = 1 \) transition, Zeeman — is located in a field of two counterpropagating linearly polarized traveling waves which is equivalent to the field of two standing waves with \( \sigma^- \) and \( \sigma^+ \) polarizations.

The total electric field can be written as

\[
E = E_0 e^{-i\omega t} (\hat{e}_x e^{i k z} + \hat{e}_y e^{-i k z}) + \text{c.c.},
\]  

(1)

where we have chosen the coordinate system in such a way (see Fig. 1) that the light waves propagating in \( -z \) and \( +z \) direction are polarized parallel to the unit vectors \( \hat{e}_x = \hat{\varepsilon}_x \cos \varphi + \hat{\varepsilon}_y \sin \varphi \) and \( \hat{e}_y \), respectively. Because of the magnetic field parallel to the light-field propagation direction (\( z \)-axis) it is very useful to rewrite the electric field in terms of its circularly polarized \( \sigma^+ \), \( \sigma^- \) components in order to calculate the atomic response. With the use of the corresponding unit vectors \( \hat{e}_z = \pm (\hat{\varepsilon}_x \pm i \hat{\varepsilon}_y) \sqrt{2} \) we obtain from equation (1):

\[
E = E_0 e^{-i\omega t} \sqrt{2} [-\hat{e}_+ e^{-i\varphi/2} \cos (kz + \varphi/2) + \hat{e}_- e^{i\varphi/2} \cos (kz - \varphi/2)] + \text{c.c.}
\]

(2)

Now it is quite easy to see that a field of two counterpropagating linearly polarized traveling waves is equivalent to the field of two standing waves, one with right-hand (\( \sigma^+ \)) and the other
with left-hand ($\sigma^-$) circular polarization. The intensities of these $\sigma^+$ and $\sigma^-$ standing-wave components are given by

$$I \pm = I_0 [1 + \cos (2 k z \pm \varphi )].$$

Thus, in the $\sigma^+$, $\sigma^-$ representation, the angle $\varphi$ between the two linear polarization vectors finds its manifestation in a phase shift $2 \varphi$ between the spatial oscillations of the $\sigma^+$ and $\sigma^-$ standing-wave intensity.

For this monochromatic laser field, let us now consider the MO dipole force exerted on the atoms with the Zeeman-split $J = 0 \rightarrow J' = 1$ transition in an analogous way as in reference [5] for two-level atoms in a bichromatic laser field. We assume that the laser frequency $\omega$ and the magnetic field $B$ are chosen in such a way that one obtains a nearly resonant interaction of the atoms with one polarization component (say, $\sigma^+$) and a strong detuning of the other component ($\sigma^-$) fulfilling the condition

$$|\Delta_-| \gg |\Delta_+|, \Omega_+(z), \Omega_-(z), \gamma .$$

Here $\Delta_\pm = \Delta_0 \pm g_j \mu_B B/\hbar$ denotes the detuning of the laser light with respect to the $\sigma^\pm$ transition frequency, and $\Omega_\pm (z) = \gamma \sqrt{2 I_\pm (z)/I_{sat}}$ is the space-dependent optical Rabi frequency associated with the $\sigma^\pm$ standing-wave component. Under condition (4), while the $\sigma^+$ transition can be strongly excited, no significant saturation occurs on the $\sigma^-$ transition. Furthermore, for $|\Delta_-| \gg |\Delta_+|$, the strong detuning $\Delta_-$ is essentially determined by the magnetic field ($\Delta \approx 2 g_j \mu_B B/\hbar$). We point out that equation (4) represents the direct analogue to the detuning condition that was used in previous work [5] to calculate the dipole force on two-level atoms in a monochromatic laser field: while there one frequency component was assumed to be tuned far out of resonance, here the Zeeman effect allows us to consider one polarization component of the monochromatic field to be far out of resonance. In both cases, the condition of one strongly detuned field component permits a relatively easy calculation of the OM force and a simple interpretation of the rectification effect.

With the use of an appropriated density matrix formalism, it is straightforward to calculate the atomic dipole moments induced on the $\sigma^+$ and $\sigma^-$ transitions. The two corresponding contributions $F_+$ and $F_-$ to the total dipole force $F_{dip} = F_+ + F_-$ are then given by

$$F_\pm = \hbar \text{Re} (p_\pm \nabla \Omega_\pm),$$

where $p_\pm$ is the slowly varying part of the off-diagonal density matrix element $\langle J = 0, m = 0 | J' = 1, m' = \pm 1 \rangle = p_\pm e^{i m' \varphi}$ in the usual rotating wave approximation. Under condition (4) considering terms up to first order in $1/\Delta_-$, and furthermore assuming sufficiently small atomic velocity components $v_\pm$ in laser beam direction ($kv_\pm \ll \gamma$), for which motion-induced effects, like, e.g., friction forces [6] play no essential role, we obtain the main dipole force contribution

$$F_+ (z) = -\frac{1}{2} \hbar \gamma \frac{\Delta_{eff}(z)/\gamma}{\Delta_{eff}(z)/\gamma^2 + 1 + G_+(z)} \nabla G_+ (z),$$

with

$$\Delta_{eff}(z) = \Delta_+ + \gamma^2 G_- (z)/2 \Delta_-$$

Here $G_\pm (z) = \Omega_\pm^2 (z)/2 \gamma^2 = I_\pm (z)/I_{sat}$ represents the saturation parameters of the $\sigma^\pm$ standing-wave components of the laser field. The effective detuning parameter $\Delta_{eff}(z)$, which characterizes the resonance behavior of the $\sigma^+$ standing wave with respect to the $m = 0 \rightarrow m' = 1$ transition, takes into account that a space-dependent light shift of the ground state is induced by the off-resonant excitation of the $m = 0 \rightarrow m' = -1$ transition by the
\( \sigma^- \) standing wave. We note that the light-shift effect here occurs by a factor 2 weaker than in the corresponding situation of a bichromatic excitation of a two-level system [5].

The light-shift-induced rectification takes place in the dipole force component \( F_+ (z) = F_+ (z) \hat{z} \). Let us consider an illustrating example, assuming \( \varphi = \pi/4 \), \( I_0 = 800 I_{\text{sat}} \), \( \Delta_- = 50 \gamma \), and \( \Delta_+ = -8 \gamma \). According to equation (6), in our example, the effective detuning \( \Delta_{\text{eff}}(z) \) spatially oscillates in a symmetrical way around the exact resonance [see Fig. 3a] as a consequence of the light shift induced by the \( \sigma^- \) standing wave. Since the MO force \( F_+ \) associated with the \( \sigma^+ \) transition occurs as an odd function of \( \Delta_{\text{eff}} \) [see Eq. (5)], this spatially alternating sign reversal of \( \Delta_{\text{eff}}(z) \) leads to a corresponding sign reversal of \( F_+ (z) \) in comparison with the case of a spatially constant detuning \( \Delta_{\text{eff}}(z) = \Delta_+ [\text{compare solid and dashed lines in Figs. 3(a, b)]}. Under optimum rectification conditions as assumed in our example, this light-shift-induced sign reversal of \( F_+ (z) \) exactly compensates for the sign reversal that would occur in \( F_+ (z) \) without the \( \sigma^- \) standing wave [see Fig. 3(b)]. This spatial rectification leads to an overall-positive force \( F_+ (z) \). The corresponding rectified force \( \langle F_+ \rangle \), defined as force average over the optical wavelength \( \lambda \), greatly exceeds the spontaneous force limit \( \hbar k \gamma \) (\( \langle F_+ \rangle = 6.4 \hbar k \gamma \) in our example).

![Fig. 3](image-url) Illustration of the light-shift-induced rectification of the dipole force component \( F_+ (z) \), assuming \( I_0 = 800 I_{\text{sat}} \), \( \Delta_- = 50 \gamma \), \( \Delta_+ = -8 \gamma \), \( \varphi = \pi/4 \). (a) solid curve — space-dependent effective detuning \( \Delta_{\text{eff}}(z) \) for \( \sigma^+ \) component of the field, dashed curve — \( \Delta_{\text{eff}}(z) \) in absence of \( \sigma^- \) component \( (I_- = 0) \). (b) solid curve — the resulting dipole force \( F_+ (z) \), dashed curve — the force in absence of the \( \sigma^- \) component.

![Fig. 4](image-url) The wavelength-averaged MO force \( \langle F_+ \rangle \) as a function of \( \varphi \) for the same parameters as assumed in figure 3.
The angle $\varphi$ obviously plays a crucial role for the rectification effect since it determines where the sign reversal of the dipole force $F_+ (z)$ takes place. Figure 4 displays a typical dependence of the total rectified dipole force or MO force $\langle F_{\text{dip}} \rangle$ on $\varphi$, which we have calculated by averaging the force resulting from equation (5) over the optical wavelength. Figure 4 demonstrates that $\langle F_{\text{dip}} \rangle$ occurs as an odd function of $\varphi$, resembling a $\sin(2 \varphi)$ dependence. For $\varphi = 0, \pm \pi/2, \pm \pi (0 \pm 90^\circ, \pm 180^\circ)$, no rectification takes place as a consequence of the unsuitable phase shift $2 \varphi = 0, \pm \pi (0^\circ, \pm 180^\circ)$ between the $\sigma^+$ and $\sigma^-$ standing wave. The optimum optical force rectification (see Fig. 3) is obtained for $\varphi = \pm \pi/4, \pm 3 \pi/4(\pm 45^\circ, \pm 135^\circ)$.

2. Classical interpretation of the MO force for an arbitrary atom.

For the basic mechanism underlying this force a simple classical interpretation can be given. We consider an atom located in a static magnetic field and irradiated by two counterpropagating traveling monochromatic linearly polarized waves as illustrated in figure 1 [7].

The absorption of a photon from traveling wave 1, in which the momentum $+ \hbar k$ is transferred to the atom along the $z$-axis, is connected with an induced electric dipole moment (IEDM) orientated along the $x$-axis. It is well known [8] that in the presence of a magnetic field $B \parallel \hat{z}$ this IEDM processes with respect to the $z$-axis with the Larmor frequency $\omega_L = eB/2m$, where a classical electron is considered with Lande factor $g = 1$. We assume that the optical Rabi frequency $\omega_R$ is chosen in such a way that the maximum probability for a stimulated emission of a photon occurs at a time $\tau = \pi/\omega_R = \pi/4 \omega_L$ after an absorption, when the IEDM has performed a $45^\circ$ rotation. At this time, the direction of the IEDM will coincide with the direction of the polarization vector of wave 2 (see Fig. 1). As a result, the probability of a stimulated emission of the photon into wave 2 will be greater than into wave 1. In total, the absorption of a photon from wave 1 and the emission into wave 2 changes the atomic momentum by $+ 2 \hbar k$.

After the absorption of a photon from wave 2 and a subsequent $45^\circ$ Larmor rotation, the IEDM will be orientated perpendicular to the polarization vector of wave 1. As a result, the probability for a stimulated emission into wave 1 (correspondingly, the atomic momentum change $- 2 \hbar k$) will be strongly decreased.

In the ideal case the minimum time for the transfer of two photons from wave 1 into wave 2 and the change of the atomic momentum by $+ 4 \hbar k$ is given by the period of the Larmor precession $2 \pi/\omega_L$. Correspondingly the force can be written in the form

$$F_{\text{MO}} = \alpha \frac{2 \hbar \omega_L}{\pi},$$

where $\alpha$ (the effectivity of the photon transfer) depends on the parameters of the laser field and the atomic transition. We point out two properties of the force: first, the force vanishes for $\varphi = n \cdot 90^\circ$ ($n = 0, 1, 2, \ldots$) as a consequence of the equal probability of stimulated photon transfers between the two waves. Second, the force changes its sign for a sign reversal of $\varphi$ or an inversion of the magnetic field direction.

3. Observation of the magneto-optical force for sodium atoms.

A scheme of our experiment to verify the existence of the magneto-optical force is shown in figure 1. A Na atomic beam was formed by two diaphragms with diameter 0.25 mm separated by 290 mm. 10 mm behind the second diaphragm the atomic beam was intersected at right angle by two counterpropagating laser beams (see Fig. 1).

The diameter of the beams was $2q = 0.4$ mm (1/e intensity drop). The laser power in each of the beams was $P = 8$ mW, which corresponds to a Rabi frequency $\omega_R = \pi \times 200$ MHz.
The polarization direction of one of the beams could be varied with the help of a lambda/2 plate. The frequency of the laser field was tuned into resonance with the transition $3^2\text{S}_{1/2}$ ($F = 2$) $\rightarrow$ $3^2\text{P}_{3/2}$ ($F' = 3$) of the Na atom.

In the interaction region of the laser radiation with the atoms a static magnetic field of $B \approx 35$ G ($\omega_L \approx 2 \pi \times 50$ MHz assuming a classical electron without considering the different $g_J$ and $g_F$ factors involved) was created using a pair of Helmholtz coils with a diameter of 1 cm. For the compensation of optical pumping which empties the level $F = 2$ we used an additional laser beam perpendicular to the other beams and tuned into resonance with the transition $F = 1 \rightarrow F' = 2$. The detection of the transverse spatial profile of the atomic beam was realized at a distance $L = 290$ mm behind the interaction region with the help of a spatially scanning laser beam. The angle of this probing laser beam with respect to the atomic beam was 78°. The frequency of the probe light was tuned + 250 MHz above the transition $F = 2 \rightarrow F' = 3$ so that only atoms with a longitudinal velocity of $v_L = 700$ m/s were detected.

Atomic beam profiles obtained for $\varphi = 45^\circ$ (see Fig. 1) are shown in figure 5. Here the frequency of the deflecting field was tuned exactly into resonance with the transition $F = 2 \rightarrow F' = 3$. As expected from the theoretical model, the direction of the deflection of the atomic beam is determined by the direction of the magnetic field. For a magnetic field in $\pm z$ direction the average displacement of the atomic beam center was $\Delta z = \pm 0.15$ mm. This corresponds to an average deflecting force $\langle F_{MO} \rangle = \frac{M_\psi^2 \Delta z}{2 q L} = \pm 0.7 \hbar k \gamma$. Thus the effectivity of the unidirectional photon transfer according to equation (7), was $\alpha = 0.11$.

In a series of experiments we measured the dependence of the atomic beam deflection on the polarization angle $\varphi$ for the direction of the magnetic field shown in figure 1. In this measurement we did not compensate for optical pumping, and the deflecting laser field was detuned by + 50 MHz with respect to the transition $F = 2 \rightarrow F' = 3$. Here, as a consequence of optical pumping only a fraction of $\sim 1/60$ of the atoms remained in the relevant $F = 2$ ground state. This measurement confirms that the magneto-optical force vanishes for parallel and perpendicular polarization of the laser beams. The maximum deflection shows up for $\varphi = \pm 45^\circ$.

![Fig. 5. — Atomic beam profiles in registration region: 1 — without deflecting laser field, 2 and 3 — with deflecting laser field and a magnetic field directed as shown in figure 1 and in opposite direction, respectively.](image-url)
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Fig. 6. — Dependence of the displacement of the atomic beam on the polarization angle $\varphi$.

Conclusion.

These experimental results verify the theoretical predictions based on a consideration of atoms with $J = 0 \rightarrow J' = 1$ transition and a classical interpretation that the magneto-optical force depends in sign on both the direction of the magnetic field and the sign of the polarization angle $\varphi$. It is furthermore confirmed that the force vanishes for parallel and perpendicular polarization of the two laser fields and exhibits maximum strength for $\varphi = \pm 45^\circ$. The maximum deflecting force $F_{\text{MO}} = 0.7 \hbar k \gamma$ was slightly below the principal maximum $\hbar k \gamma$ of the spontaneous force. The reason for this however, was no principal one but the limited laser power available in this first experiment.

References