Dark propagation modes in optical lattices

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We examine the stimulated light scattering onto the propagation modes of a dissipative optical lattice. We show that two different pump-probe configurations may lead to the excitation, via different mechanisms, of the same mode. We found that in one configuration the scattering on the propagation mode results in a resonance in the probe transmission spectrum while in the other configuration no modification of the scattering spectrum occurs, i.e. the mode is dark. A theoretical explanation of this behaviour is provided.

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I. INTRODUCTION

Brillouin scattering is the scattering of light onto a propagating acoustic wave. In spontaneous Brillouin scattering the propagating wave corresponds to thermal, or quantum-noise, fluctuations in the material medium. On the contrary, in stimulated Brillouin scattering (SBS) the density propagating wave originates from the interference pattern between a probe and an additional pump beam. The strong pump beam can then be diffracted onto the density wave in the direction of the probe, modifying in this way the probe transmission. The SBS scheme permits both the excitation of the propagation modes of a medium as well as their detection via modification of the probe transmission. It is in this way possible to determine the phonon modes of the medium, and their respective velocity.

In this work we examine the key features of the SBS process for a nonlinear medium consisting of atoms cooled in a dissipative optical lattice. This system offers significant advantages for the study of basic nonlinear optical phenomena over condensed matter samples. First, the atomic dynamics in an optical lattice is quite well understood, and can be precisely studied through Monte Carlo simulations. Second, the excitation of propagation modes in the system can be directly detected by imaging techniques. Both points are essential for the present study. We show that the same propagation mode can be excited by two different pump-probe configurations. In one case the scattering on the propagation mode results in a resonance in the probe transmission spectrum, while in the other case no modification of the spectrum occurs, i.e. the mode is in this case dark. We describe the different excitation processes of the propagation mode for the two different configurations examined, and identify the mechanism of generation of the phase-mismatch between laser fields and the material grating which inhibits the light scattering on the propagation mode.

II. THE 3D LIN⊥LIN OPTICAL LATTICE

The nonlinear medium consists of $^{85}$Rb atoms cooled and trapped in a dissipative optical lattice. These lattices are based on the Sisyphus cooling mechanism. The periodic modulation of the light polarization, produced by the interference of several laser beams, leads to a periodic modulation of the light shifts (optical potentials) of the different Zeeman sublevels of the ground state of the atom. As a result of the optical pumping between different optical potentials, atoms are cooled and finally trapped at the potential minima.

In this work we use a 3D lin⊥lin dissipative optical lattice. The arrangement of the laser fields is shown in Fig. 1: two $x$-polarized beams propagate in the $yOz$ plane and make an angle $2\theta$, and two $y$-polarized beams propagate in the $xOz$ plane and make the same angle $2\theta$. The interference pattern of the four beams create an orthorhombic potential with minima associated with pure circular (alternatively $\sigma^+$ and $\sigma^-$) light polarizat
tion. The lattice constants, i.e. the distance between two sites of equal polarization are \( \lambda_{x,y} = \lambda / \sin \theta \) and \( \lambda_z = \lambda / (2 \cos \theta) \), with \( \lambda \) the laser field wavelength. For all the measurements presented in this work the angle \( 2 \theta \) between the lattice beams is kept fixed to 60°.

The procedure to load the atoms in the optical lattice is the standard one used in previous experiments [3]. The rubidium atoms are first cooled and trapped in a magneto-optical trap (MOT). Then the MOT magnetic field and laser beams are turned off and the lattice beams are turned on. After 10 ms of thermalization of the atoms in the lattice, two additional laser fields (beams c and p of Fig. 1) are introduced for the excitation of the propagation modes. They are derived from an additional laser, with their relative detuning \( \delta = \omega_p - \omega_c \) controlled by acousto-optical modulators. These two additional laser fields are detuned with respect to the lattice beams of some tens of MHz, so that there is no atomic observable which can be excited at the beat frequency. Furthermore, as they are derived from a laser different from the one producing the lattice beams, the effect of the unwanted beat is significantly reduced. The beams c and p cross the atomic sample in the \( xOz \) plane, and they are symmetrically displaced with respect to the z axis forming an angle \( 2 \varphi \).

III. PROPAGATION MODES

A. Generalities

The propagation modes in dissipative optical lattices have been identified in Ref. [4] and shown to exhibit interesting nonlinear effects such as stochastic resonance [9, 10]. We briefly summarize their main properties. They consist of a sequence in which one half oscillation in a potential well is followed by an optical pumping process to a neighboring well, and so on (Fig. 2). In this way, the atom travels over several potential wells by regularly changing its internal state (from \( |g, +1/2 \rangle \) to \( |g, -1/2 \rangle \) in the case of a \( J_g = 1/2 \rightarrow J_e = 3/2 \) atomic transition, as considered in Fig. 2).

The velocity \( \bar{v} \) of the propagation mode is essentially determined by the intrawell dynamics. A straightforward calculation [4] shows that for a mode in the \( x \) direction this velocity is

\[
\bar{v} = \frac{\lambda \Omega_x}{2\pi \sin \theta},
\]

where \( \Omega_x \) is the \( x \) vibrational frequency at the bottom of a potential well.

B. Excitation mechanisms

The propagation modes can be excited by adding a moving potential modulation. We consider two different configurations for the modulation beams (beams c and p in Fig. 1). In both configurations the modulation beams have the same amplitude. In the first configuration, hereafter called the \( || \) configuration, both beams have \( y \) linear polarization. The light interference pattern consists of an intensity modulation moving along the \( x \) axis with phase velocity

\[
v_{mod} = \frac{\delta}{|\Delta k|} = \frac{\delta}{2k \sin \varphi},
\]

where \( \Delta k = k_p - k_c \) is the difference between the wavevectors of the modulation beams (\( |k_c| \cong |k_p| \cong k = 2\pi / \lambda \)). This configuration has already been considered in previous work [11] and it is reexamined here for comparison with the novel excitation scheme introduced in the present work. This latter, denoted as \( \perp \) configuration, consists of a \( y \) polarized beam (beam c of Fig. 1) and of a beam with linear polarization in the \( xOz \) plane (beam p). The light interference pattern consists in this case of a polarization modulation moving along the \( x \) axis with the same phase velocity \( v_{mod} \) (Eq. 2) as in the case of the \( || \) configuration.

To determine the effective excitation of the propagation modes we monitor the velocity of the center-of-mass (CM) of the atomic cloud as a function of the velocity of the applied potential modulation. This is done by direct imaging of the atomic cloud with a CCD camera. We verify that for a given detuning \( \delta \), i.e. for a given velocity \( v_{mod} \) of the moving modulation, the motion of the center of mass of the atomic cloud is uniform, and corresponding determine the CM velocity \( v_{cm} \). Experimental results for the \( x \) component \( v_{cm,x} \) of the CM-velocity as a function of \( v_{mod} \) are reported in Fig. 3 for both the \( || \) and the \( \perp \) configurations. The observed resonant behaviour of \( v_{cm,x} \) with \( v_{mod} \) is the signature of the excitation of propagation modes in the \( x \) direction. We therefore conclude that both pump-probe configurations lead to the excitation of a propagation mode in the \( x \) direction.

\[\text{FIG. 2: An atomic trajectory corresponding to a propagation mode in the } x\text{-direction. The shown potential curves (}\gplus{+} \text{ and } g_{-}\text{) are the section along } y = z = 0 \text{ of the optical potential for a } J_g = 1/2 \rightarrow J_e = 3/2 \text{ atomic transition and a 3D lin. lin beam configuration.}\]

\[\text{FIG. 3: The velocity of the atomic cloud as a function of the modulating velocity } v_{mod}. \text{ The line indicates the results for the } x\text{-component of the CM-velocity for both the } || \text{ and the } \perp \text{ configurations.}\]
FIG. 3: Top: experimental results for the $x$ component of the velocity of the center-of-mass of the atomic cloud as a function of the velocity $v_{\text{mod}}$ of the moving light interference pattern. The angle between pump and probe beams is $2\varphi = 48$. Bottom: position of the resonances as a function of the sine of the half-angle $\varphi$ between the pump and the probe beams. The points refer to experimental findings (expt.) and to semiclassical Monte Carlo simulations (MC), the lines to Eqs. (3).

To determine the nature of the observed propagation modes, we examine the atomic dynamics in the optical lattice with the help of semiclassical Monte Carlo simulations [12]. The analysis of the numerically calculated atomic trajectories shows that the excited mode is the same for both configurations, and consists of a sequence of an half oscillation in a potential well followed by an optical pumping into the neighbouring well, as in Fig. 3a.

We turn now to the analysis of the excitation mechanism of the propagation modes for the two pump-probe configuration. To this end, it is useful to examine the dependence of the position of the resonance in the velocity of the CM of the atomic cloud (as the ones in Fig. 3a) on the angle $2\varphi$ between pump and probe beams. By taking several measurements for different values of the angle $2\varphi$ between the modulation beams we determine the position of the resonances as a function of the angle $\varphi$, as reported in Fig. 3b. On the same plot we also reported results of semiclassical Monte Carlo simulations, which are found to be in very good agreement with the experimental findings. The results of Fig. 3 show that the velocity $v_{\text{mod}}$ of the light intensity interference pattern ($||$ configuration) required to excite a propagation mode differs from the velocity of the polarization grating ($\perp$ configuration) leading to the excitation of the same mode. The condition for the velocity of the light interference pattern to excite a propagation mode can be obtained by imposing that the atoms following the mode are at all times dragged by the moving potential modulation corresponding to the light interference pattern of the beams $c$ and $p$. This requires that the light polarization pattern moves at a velocity $v_{\text{mod}}$ such that the lattice potential well actually occupied by the atom gets deeper as a result of the modulation of the optical potential. The resulting conditions on the velocity of the moving modulation, and equivalently on the pump-probe detuning $\delta$, can therefore be derived by examining the effect of the modulation on the optical potentials. Consider first the $||$ configuration. The light interference pattern is a moving intensity modulation, therefore all optical potentials are modulated in phase: at a given instant and position all potential wells corresponding to the different atomic ground state Zeeman sublevels get both deeper or shallower as a result of the modulation. Thus to excite the propagation mode the atoms should follow the moving intensity modulation, i.e. the phase velocity $v_{\text{mod}}$ of the light interference pattern should be equal to the velocity $\bar{v}$ of the mode:

$$v_{\text{mod}} = \pm \bar{v} .$$

(3)

Consider now the $\perp$ configuration. The light interference pattern is in this case a moving polarization modulation, with the optical potentials associated with opposite Zeeman sublevels (opposite quantum number $m$) modulated in phase opposition. It follows that in this configuration a modulation moving at the mode velocity does not lead to the mode excitation. On the contrary, to excite the propagation mode it is necessary that the modulation moves with respect to the atoms in such a way that following the transfer of an atom from a lattice well of given circular polarization ($\sigma_+$ or $\sigma_-$) to one of opposite polarization, the modulation changes sign. Quantitatively, consider the time interval $\Delta t$ in which the atom in the propagation mode makes half an oscillation in a potential well and then is optically pumped into the neighbouring well. Then we simply have $\bar{v} \cdot \Delta t = \lambda_+/2$. In the same time interval the modulation polarization should be reversed, i.e. should change from $\sigma_+$ to $\sigma_-$ (or viceversa, depending on which potential well is initially occupied by the atom). Considering that in the time interval $\Delta t$ the atom moved of $\lambda_+/2$, and that in the moving modulation a maximum of polarization $\sigma_-$ is spaced of $\lambda_m/2 = \pi/|\Delta k|$ from the following maximum of polarization $\sigma_+$, we find that the light interference pattern should move at the velocity $v_{\text{mod}} \cdot \Delta t = \pm (\lambda_m/2 + \lambda_+/2)$. Together with $\bar{v} \cdot \Delta t = \lambda_+/2$, we find then that in the $\perp$ configuration the condition for the excitation of the propagation mode is:

$$v_{\text{mod}} = \pm \left( 1 + \frac{\sin \theta}{2 \sin \varphi} \right) \bar{v} .$$

(4)
The conditions Eqs. (3,4) are rewritten in terms of the detuning \( \delta \) as

\[
\delta = \pm \frac{2 \sin \varphi}{\sin \theta} \Omega_x \quad (| \parallel \text{ configuration} |, \quad (5)
\]

\[
\delta = \pm \left( 1 + \frac{2 \sin \varphi}{\sin \theta} \right) \Omega_x \quad (| \perp \text{ configuration} |, \quad (6)
\]

where we used Eqs. (1,2). The very good agreement (see Fig. 3) of Eqs. (5,6) with the experimental findings and with the results of semiclassical Monte Carlo simulations demonstrate the validity of our physical picture.

C. Light scattering

FIG. 4: Transmission of the probe beam as a function of the detuning between pump and probe fields, for different values of the angle \( 2\varphi \) between pump and probe beams. The lattice detuning is \( \Delta = 50 \text{ MHz} \), the intensity per lattice beam \( I_L = 5 \text{ mW/cm}^2 \). These parameters correspond to \( \Omega_x \simeq 2\pi \times 50 \text{ kHz} \) (vertical dashed lines). The left plot corresponds to the \( | \parallel \text{ pump-probe configuration} | \), the right plot to the \( | \perp \text{ one} | \).

So far we considered the effect of the light interference pattern on the atomic sample, with the excitation of propagation modes and their detection by direct imaging of the atomic cloud. The properties of the material medium can also be studied by stimulated light scattering measurements. In fact the pump-probe interference pattern may excite a material grating onto which the pump can be diffracted in the direction of the probe beam, modifying the probe transmission. That is the approach followed now, with a view to compare light scattering measurements with the previous results obtained via direct imaging of the atomic cloud.

We decreased the amplitude of one of the modulation beams (beam p) which plays now the role of the pump (or coupling) beam. We measure the probe transmission as a function of the detuning \( \delta \) for different angles between the pump and the probe beams, with results as the ones in Fig. 4. In the case of the \( | \parallel \text{ pump-probe configuration} | \), we easily identify in the probe transmission spectrum the Brillouin resonances (the resonances in Fig. 4, marked by arrows) corresponding to stimulated light scattering on the propagation modes. We verified that the dependence of the position of these resonances on the angle \( \varphi \) is in complete agreement with Eq. (6), which confirms that these resonances originate from light scattering on propagation modes. On the contrary, in the case of the \( | \perp \text{ pump-probe configuration} | \) no resonance is observed around the position corresponding to Eq. (5) (these positions are marked by filled circles in Fig. 4b). In other words, the propagation mode is dark in the \( | \perp \text{ configuration} | \).

The absence of resonances in the scattering spectrum for the propagation mode in the \( | \perp \text{ configuration} | \) can be explained by examining the phase-mismatch between the laser and the material waves. The frequency (energy) and phase-matching (momentum) conditions for the stimulated scattering process read

\[
\omega_c = \omega_p \pm \Omega , \quad (7)
\]

\[
\vec{k}_c = \vec{k}_p \pm \vec{q} . \quad (8)
\]

Here \( \vec{q} \) and \( \Omega \) are respectively the wavevector and the frequency of the light-induced material density grating and are related by the phonon dispersion relation \( \Omega = v_{\text{grating}} |\vec{q}| \), with \( v_{\text{grating}} \) the phase velocity of the moving grating. The frequency \( \Omega \) has been determined previously for both \( | \parallel \) and \( | \perp \text{ configurations} | \) (Eqs. (3,4)). As the excited mode is the same for both pump-probe configurations, the phase velocity of the material grating does not depend on the chosen configuration and is equal to the velocity \( \vec{v} \) (Eq. 4). From these values for \( \Omega \) and \( v_{\text{grating}} \) we derive, through the dispersion relation, the momentum of the material grating:

\[
|\vec{q}_\parallel| = |\Delta k| , \quad (9)
\]

\[
|\vec{q}_\perp| = |\Delta k| \left( 1 + \frac{\sin \theta}{2 \sin \varphi} \right) . \quad (10)
\]

It turns out that in the \( | \parallel \text{ configuration} | \) the momentum \( \vec{q}_\parallel = \pm |\Delta k| \vec{e}_x \) of the material grating fulfills the phase matching condition Eq. (8), and therefore the scattering on the propagation mode results in a resonance line in the probe transmission spectrum. In contrast, in the \( | \perp \text{ configuration} | \) the momentum \( \vec{q}_\perp = \pm |\vec{q}_\parallel| \vec{e}_x \) results in a phase-mismatch between the laser and the material waves. Thus, no resonance is expected in the probe transmission spectrum, in agreement with our experimental findings.

The effective creation of a moving material grating has been confirmed by semiclassical Monte Carlo simulations.
FIG. 5: Numerical results for the atomic density as a function of $x$ for the $\parallel$ (top) and $\perp$ (bottom) pump-probe configurations. The shown density distribution is stationary in a frame moving along $x$ at a velocity $\bar{v}$.

The numerical results, as the ones shown in Fig. 5, correspond to an atomic density grating moving in the $x$ direction with a velocity $\bar{v}$ for both parallel and perpendicular configurations. In the frame moving at the velocity $\bar{v}$ of the propagation mode we find a stationary modulation of the atomic density with different wavevectors for the two pump-probe configurations and in very good agreement with Eqs. (8, 9). This confirms the validity of our analysis.

IV. CONCLUSIONS

In summary, in this work we examined the stimulated light scattering onto the propagation modes of a dissipative optical lattice. Two different pump-probe configurations have been analyzed: in one the interference pattern is a modulation of the light intensity, while in the other one the pump and probe fields give rise to a modulation of the light polarization. First, we have shown that the same propagation mode is excited in the two cases, and described the two different mechanisms of excitation. Then we analyzed the light scattering on the propagation mode. Although the mode excited in the two pump-probe configurations is the same, we found that the probe transmission spectrum is completely different for the two cases. In fact, only in one configuration the mode results in a resonance in the probe transmission spectrum. For the other configuration, no trace of the mode excitation is found in the probe transmission spectrum. This behaviour was explained in terms of phase-mismatch between the laser fields and the propagating wave.

Light scattering is a powerful technique for the study of a large variety of material media. Particularly, in optical lattices it has allowed the study of local (intrawell) as well as delocalized (inter wells) dynamics. However there is not a one-to-one correspondence between the light scattering spectrum and the atomic dynamics as shown in this work where we observed and described dark propagation modes.

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