

# Atomic Recoil Effects in Slow Light Propagation<sup>1</sup>

I. Carusotto<sup>1,4</sup>, M. Artoni<sup>2</sup>, and G. C. La Rocca<sup>3,4</sup>

<sup>1</sup> Scuola Normale Superiore, 56126 Pisa, Italy

e-mail: Iacopo.Carusotto@sns.it

<sup>2</sup> INFN, European Laboratory for Non-linear Spectroscopy, 50125 Florence, Italy

<sup>3</sup> Dipartimento di Fisica, Università di Salerno, 84081 Baronissi (Sa), Italy

<sup>4</sup> INFN, Scuola Normale Superiore, 56126 Pisa, Italy

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We theoretically investigate the effect of atomic recoil on the propagation of ultraslow light pulses through a coherently driven Bose–Einstein condensed gas. For a sample at rest, the group velocity of the light pulse is the sum of the group velocity that one would observe in the absence of mechanical effects (infinite mass limit) and the velocity of the recoiling atoms (light-dragging effect). We predict that atomic recoil may give rise to a lower bound for the observable group velocities, as well as to pulse propagation at negative group velocities without appreciable absorption. © 2000 MAIK “Nauka/Interperiodica”.

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Recent experiments [1, 2] have demonstrated a reduction in the group velocity of light down to values as low as 17 m/s in coherently driven atomic samples. This was achieved by tuning the pulse frequency in the electromagnetically induced transparency (EIT) window of an optically dressed three-level atomic gas, where quantum coherence between two lower levels gives rise to a vanishing absorption along with a very steep dispersion [3]. Further improvements to the experimental setup are expected [1] to enable one to reach group velocities as small as the atomic recoil velocity. In this regime, recoil is expected to play an important role in the propagation of the pulse.

In this letter, we provide a detailed derivation of the group velocity of light pulses in a coherently driven Bose–Einstein condensed (BEC) atomic sample [4] when the effect of atomic recoil is taken into account. Apart from the well-known light-dragging effect in uniformly moving dielectrics [5], we show that the group velocity of slow light in a sample at rest under appropriate EIT conditions is given by the group velocity in the infinite mass approximation plus the velocity of the atoms which recoil following the optical process itself. Such a dragging effect imposes a lower bound to the group velocity that can be observed in typical configurations of experimental interest. For a specific level scheme and a geometry in which atoms recoil in the direction opposite to the probe wavevector, light propagation at negative group velocities without appreciable absorption is also possible. Finally, we show that the group velocity of a light pulse is not affected by atom–atom interactions at the mean-field level.

We consider a cloud of BEC atoms [4] in a three-level  $\Lambda$ -type configuration, as shown in Fig. 1. All atoms

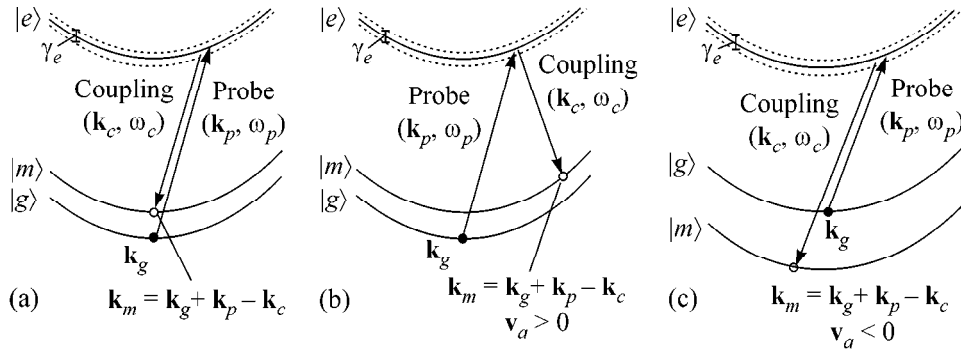
are initially in the ground state  $|g\rangle$ , and the optical transition between the metastable  $|m\rangle$  and excited state  $|e\rangle$  is dressed by a nearly resonant *coupling* cw laser beam of amplitude  $E_c(\mathbf{x})$  and frequency  $\omega_c = \omega_e - \omega_m$ . A weak *probe* pulse at frequency  $\omega_p$  nearly resonant with the other optical transition between the ground state  $|g\rangle$  and the excited state  $|e\rangle$  also propagates through the system. When the decay rate of the metastable level  $m$  is much smaller than the decay rate of the level  $e$ , the probe field experiences EIT with a narrow absorption dip and a very steep dispersion at frequencies around  $\omega_p = \omega_c + \omega_m - \omega_g$  [3]. In a second-quantized formalism, the Hamiltonian of the system can be written as

$$\mathcal{H} = \sum_{i=\{g,e,b\}} \int d^3\mathbf{x} \hat{\psi}_i^\dagger(\mathbf{x}) \left( \hbar\omega_i + V_i(\mathbf{x}) - \frac{\hbar^2}{2m} \nabla^2 \right) \hat{\psi}_i(\mathbf{x}) - (d_p E_p(\mathbf{x}, t) \hat{\psi}_e^\dagger(\mathbf{x}) \hat{\psi}_g(\mathbf{x}) + d_c E_c(\mathbf{x}, t) \hat{\psi}_e^\dagger(\mathbf{x}) \hat{\psi}_m(\mathbf{x}) + \text{h.c.}). \quad (1)$$

The first two terms describe the internal structure of the atoms and their kinetic and potential energy, while the last terms describe the coupling of the two laser beams to the atoms. The effects of the atom–atom interactions will be discussed later. Both the spontaneous emission from the excited state  $|e\rangle$  and the decoherence of the two lower  $|m\rangle$  and  $|g\rangle$  states are responsible for a loss of atoms from the condensate and can, therefore, be modeled by loss terms in the equations of motion for the three-component macroscopic wavefunction  $\Psi_i$  of the Bose condensate:

$$i\hbar \frac{\partial}{\partial t} \Psi_g(\mathbf{x}, t) = \left[ -\frac{\hbar^2 \nabla^2}{2m} + V_g(\mathbf{x}) + \hbar\omega_g \right] \Psi_g(\mathbf{x}, t) - d_p^* E_p^*(\mathbf{x}, t) \Psi_e(\mathbf{x}, t), \quad (2)$$

<sup>1</sup> This article was submitted by the authors in English.



**Fig. 1.** Level scheme and optical processes for (a) copropagating and (b) counterpropagating probe and coupling beams. (c) Proposed arrangement for obtaining negative group velocities.

$$\begin{aligned}
 i\hbar \frac{\partial}{\partial t} \Psi_e(\mathbf{x}, t) &= \left[ -\frac{\hbar^2 \nabla^2}{2m} + V_e(\mathbf{x}) + \hbar(\omega_e - i\gamma_e) \right] \Psi_e(\mathbf{x}, t) \\
 &\quad - d_p E_p(\mathbf{x}, t) \Psi_g(\mathbf{x}, t) - d_c E_c(\mathbf{x}, t) \Psi_m(\mathbf{x}, t), \quad (3) \\
 i\hbar \frac{\partial}{\partial t} \Psi_m(\mathbf{x}, t) &= \left[ -\frac{\hbar^2 \nabla^2}{2m} + V_m(\mathbf{x}) + \hbar(\omega_m - i\gamma_m) \right] \Psi_m(\mathbf{x}, t) \\
 &\quad - d_c^* E_c^*(\mathbf{x}, t) \Psi_e(\mathbf{x}, t). \quad (4)
 \end{aligned}$$

In the following, we will assume that all atoms are initially condensed in the ground state and that the probe pulse is very weak; in this case, the probe will not essentially affect the (macroscopic) condensate, so that the optical polarization caused by the noncondensed atoms generated by incoherent processes can be safely neglected. The effect of the coupling beam on the condensed atoms alone is, in fact, negligible for any value of its intensity, since its frequency is off-resonance from any optical transition starting from the ground level. For small atomic densities  $N_o(N_o/|\mathbf{k}_p|^3 \ll 1)$ , we can also assume that the photonic mode structure inside the condensed cloud is not strongly modified, compared to the free space one, so that the excited-state spontaneous emission rate  $\gamma_e$  can be taken to be the same as in free space [6]. In the spirit of a semiclassical local density approximation [4, 7], we will also neglect the effect of the external trapping potential and consider the probe and coupling beams as monochromatic plane waves of the form  $E_{p,c}(\mathbf{x}, t) = \bar{E}_{p,c} e^{i[\mathbf{k}_{p,c}\mathbf{x} - \omega_{p,c}t]}$  illuminating a locally homogeneous condensate described by the field  $\Psi_g(\mathbf{x}, t) = \bar{\Psi}_g e^{i[\mathbf{k}_g\mathbf{x} - (\omega_g + \mathbf{k}_g^2/2m)t]}$ , where  $|\bar{\Psi}_g|^2 = N_o$ . For a cloud at rest,  $\mathbf{k}_g = 0$  and  $\Psi_g(\mathbf{x}, t) = \bar{\Psi}_g e^{-i\omega_g t}$ , while for a cloud that uniformly and homogeneously moves with a velocity  $\mathbf{v}$ ,  $\mathbf{k}_g = m\mathbf{v}/\hbar$ . Due to the energy and momentum conservation, the amplitudes of the excited and metastable components

of the atomic field have the same plane-wave structure as for the ground state; i.e.,

$$\begin{aligned}
 \Psi_e(\mathbf{x}, t) &= \bar{\Psi}_e \exp\{i[(\mathbf{k}_p + \mathbf{k}_g)\mathbf{x} - (\omega_p + \omega_g^{(\text{eff})})t]\}, \quad (5) \\
 \Psi_m(\mathbf{x}, t) &= \bar{\Psi}_m \exp\{i[\mathbf{k}_p - \mathbf{k}_c + \mathbf{k}_g)\mathbf{x} - (\omega_p - \omega_c + \omega_g^{(\text{eff})})t]\}. \quad (6)
 \end{aligned}$$

Inserting these forms into Eq. (4) and then into Eq. (3) yields

$$\bar{\Psi}_m = \frac{-d_c^* \bar{E}_c^*}{\hbar(\Delta_m(\mathbf{k}_p, \omega_p) + i\gamma_m)} \bar{\Psi}_e \quad (7)$$

and

$$\bar{\Psi}_e = \frac{-d_p \bar{E}_p}{\hbar(\Delta_e(\mathbf{k}_p, \omega_p) + i\gamma_e - |d_c E_c|^2/(\Delta_m(\mathbf{k}_p, \omega_p) + i\gamma_m))} \bar{\Psi}_g, \quad (8)$$

which generalizes the expression used for describing EIT in the  $\Lambda$ -type three-level atomic configuration by including kinetic-energy corrections associated with the atomic recoil. These appear in the detuning from the excited level

$$\Delta_e(\mathbf{k}_p, \omega_p) = \omega_g^{(\text{eff})} + \omega_p - \omega_e^{(\text{eff})}(\mathbf{k}_p) \quad (9)$$

and from the metastable level

$$\Delta_m(\mathbf{k}_p, \omega_p) = \omega_g^{(\text{eff})} + \omega_p - \omega_c - \omega_m^{(\text{eff})}(\mathbf{k}_p), \quad (10)$$

where  $\omega_g^{(\text{eff})} = \omega_g^{(\text{eff})}(\mathbf{k}_g) = \omega_g + \hbar\mathbf{k}_g^2/2m$ ,  $\omega_e^{(\text{eff})}(\mathbf{k}_p) = \omega_e + \hbar(\mathbf{k}_p + \mathbf{k}_g)^2/2m$ , and  $\omega_m^{(\text{eff})}(\mathbf{k}_p) = \omega_m + \hbar(\mathbf{k}_p - \mathbf{k}_c + \mathbf{k}_g)^2/2m$ . Only the dependence on  $\mathbf{k}_p$  and  $\omega_p$ , which will be needed in the following, is explicitly indicated, whereas the dependence on the other setup parameters  $\omega_c$ ,  $\mathbf{k}_c$ , and  $\mathbf{k}_g$  is left implicit. Since the dipole moment per unit volume at the probe frequency is given by  $d_p^* \bar{\Psi}_g^* \bar{\Psi}_e$ , Eq. (8) leads to a simple expression for the

dielectric function  $\epsilon(\omega_p, \mathbf{k}_p)$  of the dressed atomic cloud,

$$\epsilon(\omega_p, \mathbf{k}_p) = 1 + \frac{4\pi N_o |d_p|^2}{\hbar(|\Omega_c|^2 / (\Delta_m + i\gamma_m) - \Delta_e - i\gamma_e)}, \quad (11)$$

where  $\Omega_c = |d_c E_c|/\hbar$  is the Rabi frequency of the coupling beam. If the spontaneous decay rate  $\gamma_e$  is much larger than all other frequency scales and, in particular, if  $\gamma_e \gg \Delta_e$ , then the detuning  $\Delta_e$  of the excited state can be neglected in Eq. (11). If we further assume that the decoherence rate  $\gamma_m$  is much smaller than  $\Gamma = \Omega_c^2/\gamma_e$ , then Eq. (11) simplifies to

$$\epsilon(\omega_p, \mathbf{k}_p) = 1 + \frac{4\pi N_o |d_p|^2}{\hbar\gamma_e} \left\{ i + \frac{\Gamma}{\Delta_m(\mathbf{k}_p, \omega_p) + i\Gamma} \right\}. \quad (12)$$

Providing the Rabi frequency  $\Omega_c$  of the coupling beam is smaller than the excited state linewidth  $\gamma_e$ , nearly total transmission occurs within a small bandwidth  $\Gamma$  of frequencies around  $\omega_p^{(o)} = \omega_m^{(\text{eff})}(\mathbf{k}_p^{(o)}) + \omega_c - \omega_g^{(\text{eff})}$ , for which  $\Delta_m(\mathbf{k}_p^{(o)}, \omega_p^{(o)}) = 0$ ; in the same frequency window, the refractive index, which is unity ( $\omega_p^{(o)} = c|\mathbf{k}_p^{(o)}|$ ) at line center, has a very steep dispersion. This implies that a narrow-band pulse would propagate with a very small group velocity without being appreciably absorbed [1, 2, 8]. Approximating the atomic dispersion of the metastable  $|m\rangle$  state after the absorption of a photon from the probe beam and its immediate reemission into the coupling beam as a linear one with the group velocity

$$\mathbf{v}_a = \mathbf{v} + \frac{\hbar}{m}(\mathbf{k}_p^{(o)} - \mathbf{k}_c), \quad (13)$$

the detuning in the denominator of Eq. (12) can be approximated by  $\Delta_m(\mathbf{k}_p, \omega_p) \simeq (\omega_p - \omega_p^{(o)}) - (\mathbf{k}_p - \mathbf{k}_p^{(o)})\mathbf{v}_a$ , so that  $\epsilon(\omega_p, \mathbf{k}_p)$  acquires the new form

$$\epsilon(\omega_p, \mathbf{k}_p) \simeq 1 + \frac{4\pi N_o |d_p|^2}{\hbar\gamma_e\Gamma} [(\omega_p - \omega_p^{(o)}) - (\mathbf{k}_p - \mathbf{k}_p^{(o)})\mathbf{v}_a]. \quad (14)$$

The dispersion law for a probe propagating in the direction of the unit vector  $\hat{\mathbf{k}}_p = \mathbf{k}_p/|\mathbf{k}_p|$  with a frequency centered on the EIT transparency window can be obtained by inserting Eq. (14) into  $\epsilon(\omega, \mathbf{k})\omega^2 = c^2\mathbf{k}^2$  and then linearizing around  $\omega_p = \omega_p^{(o)}$  and  $\mathbf{k}_p = \mathbf{k}_p^{(o)}$ :

$$(\omega_p - \omega_p^{(o)}) = (\mathbf{k}_p - \mathbf{k}_p^{(o)}) \frac{\eta \mathbf{v}_a + c\hat{\mathbf{k}}_p}{1 + \eta}, \quad (15)$$

where  $\eta = 2\pi N_o |d_p|^2 \omega_p^{(o)} / \hbar \Omega_c^2$ . The relevant group velocity  $\mathbf{v}_g = \nabla_{\mathbf{k}_p} \omega_p$  at two-photon resonance can finally be written as

$$\mathbf{v}_g = \frac{c}{1 + \eta} \hat{\mathbf{k}}_p + \frac{\eta}{1 + \eta} \mathbf{v}_a. \quad (16)$$

For a sample at rest, in the infinite mass limit,  $\mathbf{v}_a$  is negligible and the group velocity has the usual expression  $\mathbf{v}_g = c\hat{\mathbf{k}}_p/(1 + \eta)$  [8]. In this case, for values of  $\eta$  much larger than unity, light speeds much less than  $c$  can be observed as, e.g., in [1], where  $\eta \sim 10^7$ . However, we cannot neglect atomic recoil when  $\eta$  is much larger than unity and is of the order of  $c/|\mathbf{v}_a|$ , since  $\mathbf{v}_g$  becomes comparable in magnitude to  $\mathbf{v}_a$ . In this case, the group velocity can be written as

$$\mathbf{v}_g \simeq \frac{c}{\eta} \hat{\mathbf{k}}_p + \mathbf{v}_a. \quad (17)$$

While the first term  $c\hat{\mathbf{k}}_p/\eta$  recovers Eq. (1) in [1], the other term seems to suggest that light is dragged by the metastable atoms, which recoil at a velocity of  $\mathbf{v}_a$ ; however, we stress that, under our conditions,  $|\bar{\Psi}_e|^2 + |\bar{\Psi}_m|^2 \ll |\bar{\Psi}_g|^2 = N_o$  and, therefore, the center-of-mass motion of the atomic cloud is weakly affected by light.

We now proceed to discuss novel and interesting effects associated with result (17). For an atomic sample at rest, in which  $|g\rangle$  and  $|m\rangle$  are hyperfine sublevels of the same ground state with energies very close to each other,  $\mathbf{k}_g = 0$  and  $\mathbf{v}_a$  turns out to be a negligibly small quantity for copropagating probe and coupling beams (Fig. 1a). Such a situation was examined in [7], e.g., where recoil is explicitly omitted. On the other hand, for counterpropagating beams (Fig. 1b),  $\mathbf{v}_a$  is nearly twice the recoil velocity of the  $|g\rangle \rightarrow |e\rangle$  optical transition and it is directed as the probe wavevector; in such a geometry, the group velocities are then restricted by the lower bound  $|\mathbf{v}_a|$ . In the case of sodium atoms ( $D_2$  line), this quantity is approximately 6 cm/s, i.e., 300 times smaller than the lowest group velocity of 17 m/s so far reported in sodium [1]. Since the most stringent upper bound to  $\eta$  is actually set by the lower bound to the coupling intensities  $\Omega_c^2 > \gamma_m\gamma_e$ , which have to be applied in order for the EIT to be fully developed, a substantial reduction of  $\gamma_m$  [1] will lead to much larger values of  $\eta$ , so that the effect of atomic recoil, as predicted by Eq. (17), could possibly be observed.

For a sample moving with the uniform velocity  $\mathbf{v}$ , our theory recovers the well-known Fresnel–Fizeau light-drag [5] effect; in the slow-light case, all velocities involved are nonrelativistic and the Galilean composition of velocities is obtained as in Eq. (17). Unlike the effect of atomic recoil, Fresnel–Fizeau drag occurs even in the infinite atomic mass limit. Recently, a related effect was shown to lead to exotic features of

light propagation in the more complex situation of non-uniformly moving media [9], but this is beyond the scope of this paper.

With copropagating coupling and probe beams and the appropriate choice of the atomic levels, i.e., the  $\Lambda$  configuration, in which the level  $m$  has an energy lower than  $g$  (see Fig.1c), the recoil velocity  $\mathbf{v}_a$  is directed oppositely with respect to the probe beam even for a sample initially at rest. In this case, for sufficiently small values of  $c/\eta$ , the probe wavevector and the group velocity turn out to be oppositely directed. From a phenomenological point of view, the possibility of attaining such *negative* group velocities may be exploited to investigate rather novel effects in the domain of geometrical optics, such as, e.g., negative refraction angles at the boundary with free space [10]. Recent developments in coherently prepared atomic media have revived the interest in the issue of negative group velocities. With respect to the previous works on the subject [11, 12], our proposal is characterized by the fact that both absorption and group velocity dispersion almost vanish in the frequency range of interest, so that the shape of the light pulse remains essentially unchanged. Negative group velocities were also predicted to occur in an EIT configuration for coupling and probe beams copropagating in a *hot* atomic gas [13]: because of the Doppler effect, light interacts only with a narrow class of atomic velocities and the sample behaves as an effectively moving one. If the selected atoms move in the opposite direction with respect to the probe wavevector, negative group velocities may occur for sufficiently dense samples, as is also predicted by the present treatment when a nonzero atomic velocity is explicitly included in Eq. (13).

In actual experiments, a nonzero temperature and the finite size of the sample may cause a finite velocity spread for the ground-state atoms. This can be taken into account by integrating dielectric susceptibility (12) over the velocity distribution of ground-state atoms. For a Lorentzian velocity distribution [13], a straightforward calculation leads to the same form of susceptibility, where  $\Gamma$  in the denominator is replaced by  $(\Gamma + \Gamma_D)$ ,  $\Gamma_D = |\mathbf{k}_p - \mathbf{k}_c|v_D$  being the Doppler width expressed in terms of the velocity spread  $v_D$ . In physical terms, the effect of a Doppler width  $\Gamma_D$  comparable to the subnatural linewidth  $\Gamma$  is similar to the effect of having a lower level decoherence  $\gamma_m$  of the order of  $\Gamma$ , i.e., a broadened absorption dip and a reduced contrast of the transparency feature, which is no longer complete. From a quantitative point of view, the broadening due to the finite size of a zero-temperature BEC is generally smaller than the recoil velocity and thus can be safely neglected with respect to  $\Gamma$ . For hot samples,  $\Gamma_D$  is negligible only if  $\mathbf{k}_p \approx \mathbf{k}_c$ , i.e., for copropagating coupling and probe beams and small lower state energy splitting. In addition, if the Doppler broadening  $|\mathbf{k}_p|v_D$

of the excited state is comparable to its linewidth  $\gamma_e$ , the detuning  $\Delta_e$  can no longer be neglected in Eq. (11) and a more detailed treatment has to be carried out [13].

The theory described up to now neglected the atom–atom interactions (collisions). These are commonly modeled [4] by adding quartic terms to Hamiltonian (1) and give rise to additional cubic terms of the form  $\sum_j G_{i,j} |\psi_j|^2 \psi_i$  in the mean-field wave Eqs. (2)–(4). The coupling coefficients  $G_{i,j}$  are proportional to the  $s$ -wave scattering length for collisions between atoms in the  $i$  and  $j$  states ( $i, j = \{e, g, m\}$ ), respectively. To the lowest order in the probe intensity, only the  $G_{i,g} |\psi_g|^2 \psi_i$  terms contribute to the sum causing a mean-field shift of the  $e$  and  $m$  level frequencies in Eqs. (9) and (10). The excited level  $e$  is adiabatically eliminated in the present treatment, while the collisional frequency shift of the metastable level  $m$  gives rise to a small shift of the two-photon resonance condition in Eq. (12). This means that the photon dragging effects of interest originate from the independent recoil of each atom and, thus, Bogoliubov’s sound velocity  $v_s = \sqrt{G_{g,g} N_o/m}$  does not appear to be relevant to the linear propagation of light pulses in condensed media under EIT. The dispersion of Bogoliubov’s phonons [4] may, on the other hand, be crucial in more complex optical processes which involve the excitation of phonons in the condensate, such as, e.g., Brillouin scattering by density fluctuations [14].

In conclusion, we have shown that even in a sample at rest, under appropriate EIT conditions, light can be dragged by the atoms which recoil after the absorption of a photon from the probe beam and the subsequent emission into the coupling beam. We hope that a feasible upgrade of the experimental setup commonly used to study light propagation in EIT configurations [1] will soon allow the detection of such atomic recoil effects.

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