

Light Transport in Cold Atoms and Thermal Decoherence

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We use the coherent backscattering interference effect to investigate experimentally and theoretically how coherent transport of light inside a cold atomic vapor is affected by the residual motion of atomic scatterers. As the temperature of the atomic cloud increases, the interference contrast decreases dramatically. This emphasizes the role of motion-induced decoherence for resonant scatterers even in the sub-Doppler regime of temperature. We derive analytical expressions for the corresponding coherence time.

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Originally motivated by astrophysical purposes, wave transport in opaque media was first analyzed by means of a detailed balance of energy transfers within the scattering medium (radiative transfer theory) [1]. This energy book-keeping approach essentially relies on a random-phase assumption which discards all possible interference effects inside the medium. In its simplest form, this theory predicts a diffusive transport with spatial diffusion constant $D = \ell^2/3\tau$, where ℓ is the scattering mean free path and τ the transport time. However, it is now clear that interference can survive the average over disorder and dramatically alter the wave transport. This leads to the physics of weak and strong localization [2]. There has thus been much effort in solid state physics to study, master, and circumvent any possible phase-breaking mechanism to fully access the regime of coherent transport [3]. Coherence loss phenomena can be characterized by a phase-breaking time τ_ϕ and a corresponding coherence length $L_\phi = \sqrt{D\tau_\phi}$ [2] beyond which interference effects are essentially washed out. The coherent transport regime (also known as the mesoscopic regime) is then reached when the medium size L is smaller than L_ϕ . For instance, in metals or semiconductors, it prevails up to $L_\phi \approx 1\text{--}10 \mu\text{m}$ even at very low temperatures and very pure samples. This explains why devices with micrometer size are usually required to study mesoscopic physics [2].

During the last two decades, the field got a renewed attention in systems where electromagnetic waves were used [4]. Among the large variety of scattering media employed so far, cold atomic vapors have recently emerged. An important feature of these media is the sharp frequency resonance of the scattering cross section (of order λ^2 at resonance, λ being the wavelength). This allows for a continuous tuning of the scattering mean free path ℓ via the light frequency. Phase-breaking mechanisms in cold atomic vapors can be efficiently probed with coherent backscattering (CBS), a paradigmatic two-wave interference effect in multiple scattering [5]. In the limit of

low light intensity, an important mechanism has been identified [6]. It is rooted in the Zeeman degeneracies of the atomic internal structure but could be healed with an external magnetic field [7]. Another important phase-breaking mechanism is the residual thermal motion of atomic scatterers inside the cold vapor. This is the main topic of this Letter. Indeed, because of the Doppler and recoil effects, light frequency is changed at scattering and transport is no longer elastic. In previous papers [8,9], we have shown how the incoherent transport is affected by such effects. For typical alkali magneto-optical traps (MOT), Doppler-induced frequency shift is the dominant effect. Not surprisingly, an important parameter is the ratio of the typical Doppler shift kv (where v is the 1D rms velocity, and $k = 2\pi/\lambda$ the incoming light wave vector) to the width Γ of the atomic resonance. When $kv \gg \Gamma$, a single scattering event is enough to bring the photon completely out of resonance and the transport of light is significantly altered [10]. When $kv \ll \Gamma$, a single scattering event only slightly modifies the photon frequency and leads, as a cumulative effect, to a diffusion in frequency space [9].

In this Letter, we report the first unambiguous experimental evidence of a phase-breaking mechanism in light transport induced by the residual atomic motion [11]. We first describe the experimental procedure and our main result. We then give an analysis of the physical ingredients at the heart of the light coherence reduction. We finally compare our data to the Monte Carlo simulations described in [12], appropriately modified to include the effect of the atomic velocity distribution [13].

Our CBS experimental setup has been described in detail elsewhere [12]. The main difference here is the heating procedure to increase v . For this purpose the magnetic gradient of the MOT is switched off and the cloud is then exposed to 200 μs long, slightly red-detuned, optical molasses. An increasing amount of heating is obtained for molasses frequencies closer to resonance. To get the re-

sulting 3D velocity distribution, atoms are released from the trap and their ballistic expansion is recorded. Even though this heating technique can lead to non-Gaussian velocity distributions, we define a temperature of the cloud by $k_B T = mv^2$. The number of atoms in the cloud can be adjusted independently of T so as to maintain a fixed optical thickness $b = L/\ell$. In Fig. 1 we observe a fast decrease of the CBS enhancement factor as v is increased. The value obtained for $v \rightarrow 0$ is set by the rubidium Zeeman degeneracy. Note that the decay is faster than expected from the naïve criterion $kv \approx \Gamma$. The measurements are performed in the helicity nonpreserving channel with a resonant laser (detuning $\delta = \omega - \omega_0 = 0$, where ω and ω_0 are the laser and atomic angular frequencies). The measured angular width of the CBS peak $\Delta\theta \approx 1/k\ell$ ($k\ell \approx 1000$ typically) varied by less than 10% in the course of the experiment. This proves that the reported decrease is not due to the angular resolution of the apparatus. Before we comment on Fig. 1, let us discuss the physical ingredients of motion-induced CBS reduction for the simple double scattering configuration shown in Fig. 2. We assume here that the atomic motion is classical [14]. Atoms, distributed in space with number density $\rho \ll k^3$, are illuminated by a plane wave of angular frequency ω and wave vector \mathbf{k} . They constitute a dilute effective medium characterized by a complex index of refraction $n \approx 1 + \rho\alpha/2$, convolved by the atomic velocity distribution. The atomic complex polarizability is $\alpha = -3\pi\Gamma k^{-3}/(\delta + i\Gamma/2)$. The CBS signal can be understood as an interference between a “direct” scattering path (solid arrows) and its “reverse” counterpart (dashed arrows). To each path is associated a complex amplitude incorporating both scattering and

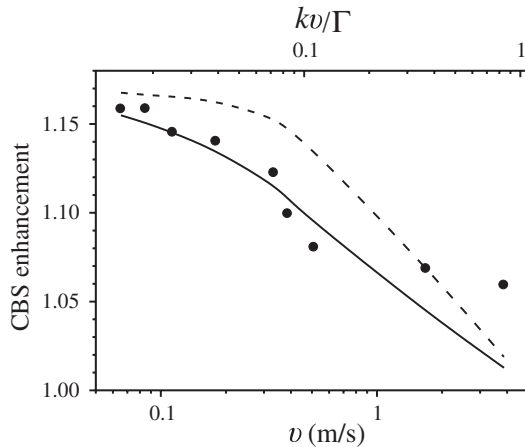


FIG. 1. CBS enhancement factor (●) as a function of the 1D typical atomic velocity v and its comparison to a Monte-Carlo calculation assuming either a Gaussian (dashed line) or a Lorentzian (solid line) atomic velocity distribution. The CBS peak is recorded in the helicity nonpreserving channel, with resonant laser light ($\delta = 0$). The optical thickness of the cloud is set constant at $b \simeq 13$. For rubidium, the velocity scale Γ/k is 4.6 m/s.

propagation in the effective medium. For scatterers at rest, with no Zeeman degeneracy, these amplitudes are exactly balanced at backscattering (outgoing wave vector $-\mathbf{k}$) in the parallel polarization channels. This yields full CBS contrast. The reduced contrast obtained for Zeeman degenerate atoms, such as rubidium, has been described in [6]. Let us now consider the impact of the atomic motion. Because of the Doppler effect, the light frequency in the laboratory reference frame is modified. For the direct path, it is $\omega_d = \omega + (\mathbf{k}' - \mathbf{k}) \cdot \mathbf{v}_1$ where \mathbf{k}' is the intermediate wave vector. After scattering by atom 2, the outgoing frequency is $\omega' = \omega + (\mathbf{k}' - \mathbf{k}) \cdot \mathbf{v}_1 - (\mathbf{k} + \mathbf{k}') \cdot \mathbf{v}_2$. For the reverse path, the frequency between the two atoms is $\omega_r = \omega - (\mathbf{k}' + \mathbf{k}) \cdot \mathbf{v}_2 \neq \omega_d$, while the final frequency is again ω' . Thus, although the outgoing waves have identical frequencies and do interfere, the intermediate frequencies differ by an amount $\simeq kv$. The CBS contrast will now be reduced as different complex amplitudes interfere. Furthermore, the frequencies seen by the atoms in their rest frames are different along the direct and reverse paths. This leads to different scattering phase shifts. One can estimate the total phase difference $\Delta\Phi$ induced by the atomic motion between the direct and reverse paths. For small velocities $kv \ll \Gamma$, one gets:

$$\Delta\Phi = \Delta\Phi_s + \Delta\Phi_p \simeq [\partial_\omega \Phi_s + k\ell \partial_\omega n]kv \approx kv/\Gamma. \quad (1)$$

The first term corresponds to the scattering phase shift, while the second describes the phase shift associated to the propagation over a distance ℓ in the effective medium. This rough estimation provides useful insights about the impact of δ at finite v (see below). For larger scattering orders ($N > 2$), the elementary scattering and propagation processes have to be chained and randomly distributed phase differences will accumulate. This decreases further the

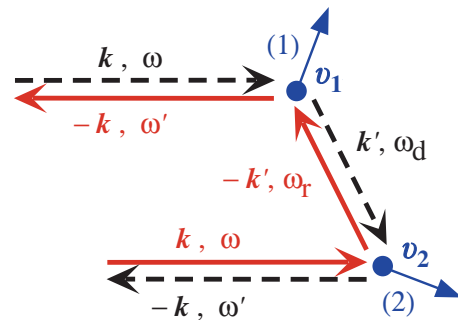


FIG. 2 (color online). CBS dephasing induced by the atomic motion, in the simplest case of two atoms moving with different velocities \mathbf{v}_1 and \mathbf{v}_2 . The CBS effect builds on the two-wave interference between amplitudes associated to reversed multiple scattering paths (arrows). The Doppler effect induces frequency redistribution of the scattered light. Although the outgoing frequencies are identical, the intermediate frequencies differ, yielding a phase-coherence loss and thus a reduction of the CBS contrast.

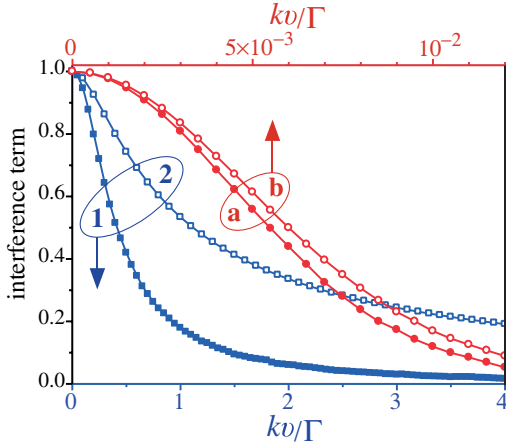


FIG. 3 (color online). Numerical calculation of the CBS interference decay for multiple scattering order $N = 2$ (\square) and $N = 40$ (\circ). The upper velocity range refers to curves **a** and **b**, while the lower one refers to curves **1** and **2**. Note the very different velocity ranges. We compare two situations: resonant excitation $\delta = 0$ (**1** and **a**) and detuned excitation $\delta \gg kv$ ($\delta = 50\Gamma$ for **2** and $\delta = 5\Gamma$ for **b**). Although the CBS contrast is improved, a detuned excitation is not sufficient to circumvent the motion-induced decoherence.

CBS contrast, see Fig. 3. An accurate computation of the shape and height of the CBS peak must include many features. First, one must average over all atomic degrees of freedom, both external (position and velocity) and internal (magnetic quantum number). Second, one must also incorporate the experimental geometry such as the shape, size, optical thickness of the cloud, as well as the waist and linewidth of the CBS probe laser beam. A Monte Carlo approach as described in [12] is then necessary, where we now include the effect of the atomic velocity distribution [8]. As mentioned earlier, some of the measured velocity distributions exhibited clear deviations from a Gaussian with an excess of fast atoms. Such atoms are responsible for large Doppler effects and are thus likely to kill the CBS interference more efficiently. We have used in our simulations either a purely Gaussian velocity distribution or a Lorentz-like distribution [where we define $v = v_0/(\sqrt{8\ln 2})$, v_0 being the FWHM velocity]. The results, displayed in Fig. 1, show a good qualitative agreement, with a fast decay of the enhancement factor with increasing v . However, the Gaussian distribution shows deviations at low v while the Lorentz distribution gives a better agreement. This pinpoints the important role played by a small fraction of atoms with large velocities. A quantitative agreement would thus require a better knowledge and control of the initial atomic velocity distribution. As can be seen, the CBS contrast decreases faster than naively expected. This is because higher scattering orders contribute significantly to the CBS effect. Indeed, at large optical thickness ($b > 10$), they contribute to a combined $\approx 30\%$ of the CBS peak height. As increasing scattering orders are

more and more sensitive to motion-induced phase breaking, this can explain our experimental observations.

As the scrambling of CBS has its origin in the frequency-dependent response of the sample (which is maximum around resonance), one could expect to recover a full interference contrast by detuning the laser to $|\delta| \gg kv$. This is an important issue for the prospective observation of the CBS effect with room-temperature vapors. Figure 3 displays Monte-Carlo simulations of the CBS contrast as a function of the rms velocity v at $\delta = 0$ and $\delta \gg kv$. Calculations have been made for $N = 2$ (squares) and $N = 40$ (circles) in a semi-infinite medium. As can be seen (curves **a** and **b**, $N = 40$), at low velocities $kv \ll \Gamma$ the interference contrast is nearly frequency independent. This behavior is due to a compensation of the variations of the scattering and propagation terms. Indeed, the phase difference in Eq. (1) is essentially proportional to the transport time τ (bracketed term), which was shown to be frequency independent near the resonance and equal to Γ^{-1} [8]. As mentioned earlier and expressed in Eq. (2), the decay rate is faster for higher scattering orders. At larger velocity (curves **1** and **2**, $N = 2$), Fig. 3 shows an increased frequency dependence of the interference contrast, which we interpret as an unperfect balance between the scattering and propagation phase shifts in Eq. (1). The interference contrast is partially restored in the $|\delta| \gg kv$ limit, where the motion-induced decoherence is now essentially driven by the propagation term in Eq. (1) of order kv/Γ . The interference thus still decreases towards zero with increasing v .

Beyond the specific case of CBS, one can quantitatively characterize the motion-induced phase breaking by defining the phase-breaking time τ_ϕ and its associated coherence length L_ϕ . We discuss here the simplest low-velocity case, $kv \ll \Gamma$. At each scattering event, the photon detuning is slightly modified and this cumulative process turns into a random walk at sufficiently large scattering orders $N \gg 1$. The width of its frequency distribution is $\approx kv\sqrt{N}$ which we also assume to be $\ll \Gamma$. If the atomic velocities have a Gaussian distribution and are independent, the average over the degrees of freedom can be performed analytically, and the resulting interference contrast decreases like:

$$c_N(v) = \exp\left[-\frac{N^3}{3}\left(\frac{kv}{\Gamma}\right)^2\right]. \quad (2)$$

Note that this quantity is independent of the initial detuning δ as discussed before. This expression is also valid for larger velocities, in the off-resonant limit $\delta \gg \sqrt{N}kv$ where the convolution by the velocity distribution can be neglected. The decay of the interference contrast described by Eq. (2) is not exponential with N , in contrast to other sources of decoherence like the atomic internal structure [6]. Nevertheless, we define a critical scattering order $N_\phi = 3(3kv/\Gamma)^{-2/3}$. The motion-induced phase-breaking

mechanism is thus characterized by:

$$\tau_{\Phi} = 3 \left(\frac{3kv}{\Gamma} \right)^{-2/3} \Gamma^{-1} \quad (3a)$$

$$L_{\Phi} = \left(\frac{3kv}{\Gamma} \right)^{-1/3} \ell. \quad (3b)$$

Therefore, the mesoscopic condition $L_{\Phi} > L$ reads:

$$\frac{kv}{\Gamma} < \frac{1}{b^3}. \quad (4)$$

For our lowest temperature of $\approx 40 \mu\text{K}$ ($v = 6.5 \text{ cm/s}$), Eq. (3b) yields $L_{\Phi} \approx 3\ell$ and condition (4) is fulfilled for optical thicknesses $b \leq 3$.

At this point, an interesting comparison with electronic transport in metals can be made. There, the temperature dependence of τ_{Φ} is governed by several different mechanisms: electron-electron scattering ($\propto T^{-2/3}$), electron-phonon scattering ($\propto T^{-3}$), and spin-flip scattering (weakly T dependent) which dominates at low T [3,15]. In our case, τ_{Φ} is set by the internal structure [6] (independent of T) and by the motion-induced decoherence, which, according to Eq. (3a), yields $\tau_{\Phi} \propto T^{-1/3}$. At $T = 40 \mu\text{K}$, τ_{Φ} is mostly determined by the internal structure (see Fig. 1).

A simple dynamical picture of the motion-induced coherence loss has been given by Golubentsev [16]. In this picture, phase coherence is lost when the length of a multiple scattering path varies (due to the motion of the scatterers) by λ during the time taken by the light to follow this path. Golubentsev found expressions equivalent to ours, with, however, an important difference: for the non-resonant case he considered, the transport time is $\tau = \ell/c < 1 \text{ ps}$ while for our narrow resonance this time is $\Gamma^{-1} \approx 27 \text{ ns}$, more than 4 orders of magnitude higher. It is only because photons are “slowed down” inside the resonant medium that we are able to observe the destruction of phase coherence.

In conclusion, we reported the first observation of a motion-induced phase breaking of the CBS effect. The observed CBS reduction is due to the fast-varying frequency response of the medium in the vicinity of a resonance, which enhances the impact of frequency redistribution caused by Doppler effect. The rapid decrease of the interference contrast indicates that the phase-coherence length L_{Φ} associated with the residual motion of atoms in our coldest sample is not very large compared to the scattering mean free path, and smaller than the actual size of the cloud. According to the mesoscopic criterion derived in this Letter, extremely low temperatures, in the nK range, are probably required to fully access interference effects in optically thick cold atomic gases. In this regime,

however, the recoil effect is expected to become the main source of decoherence.

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