Suppression of Transport of an Interacting Elongated Bose-Einstein Condensate in a Random Potential

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D. Clément,[1] A.F. Varón,[1] M. Hugbart,[1] J.A. Retter,[1] P. Bouyer,[1]
L. Sanchez-Palencia,[1] D.M. Gangardt,[2] G.V. Shlyapnikov,[2,3] and A. Aspect[1]

[1] *Laboratoire Charles Fabry, Institut d'Optique, Université Paris-Sud XI, 91403 Orsay cedex, France*
[2] *Laboratoire de Physique Théorique et Modèles Statistiques, Université Paris-Sud XI, 91405 Orsay cedex, France*
[3] *Van der Waals-Zeeman Institute, University of Amsterdam, Valckenierstraat 65/67, 1018 XE Amsterdam, The Netherlands*

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We observe the suppression of the 1D transport of an interacting elongated Bose-Einstein condensate in a random potential with a standard deviation small compared to the typical energy per atom, dominated by the interaction energy. Numerical solutions of the Gross-Pitaevskii equation reproduce well our observations. We propose a scenario for disorder-induced trapping of the condensate in agreement with our observations.

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Atomic Bose-Einstein condensates (BEC) placed in an optical potential constitute a remarkable system in which one can revisit standard problems of condensed matter physics, *e.g.* superfluidity and quantum vortices, the superfluid to Mott insulator transition, or Josephson arrays [1]. Another important topic in condensed matter physics is the transport properties in disordered materials. This is a difficult problem and it has led to the introduction of intriguing and non-intuitive concepts, *e.g.* Anderson localization [2, 3], percolation dynamics [4] and Bose [5] and spin [6] glasses. These issues have relevance to normal metallic conduction, superconductivity and superfluid flow in low temperature quantum liquids. It has a counterpart in wave physics, *e.g.* in optics and acoustics, with special interest in coherent diffusion in inhomogeneous media [7]. The main difficulty in understanding quantum transport arises from the subtle interplay of interference, scattering onto the potential landscape, and (whenever present) interparticle interactions.

Transport properties of BECs in periodic optical lattices have been widely investigated, showing lattice-induced reduction of mobility [8, 9, 10] and self-trapping [11]. Within the context of random potentials, most of the recent theoretical efforts have considered disordered or quasi-disordered optical lattices where a large variety of phenomena have been discussed such as the Bose-glass phase transition [12], localization [12, 13], and the formation of Fermi-glass, quantum percolating and spin-glass phases in Fermi-Bose mixtures [14]. Effects of disorder on BECs have also been addressed in connection to superfluid flows in liquid helium in porous media [15]. In particular, the depletions of the condensate and of the superfluid fractions have been calculated in [16], and significant shift and damping of sound waves have been predicted in [17]. Apart from the (undesired) fragmentation effect of a rough potential on trapped cold atoms and BEC's on atom chips [18], there are few experiments on BECs in random potentials [19].

In this Letter we report on the strong reduction of mobility of atoms in an elongated Bose-Einstein condensate in a random potential [20]. Starting from a quasi-1D BEC in a 3D elongated harmonic trap, we turn off the axial trapping potential while maintaining strong transverse confinement, and we monitor both (i) the axial expansion driven by the repulsive interactions [21] and (ii) the motion of the center of mass of the condensate. When the BEC is subjected to a 1D random potential created by a laser speckle, the axial expansion is strongly inhibited and the condensate eventually stops expanding (see

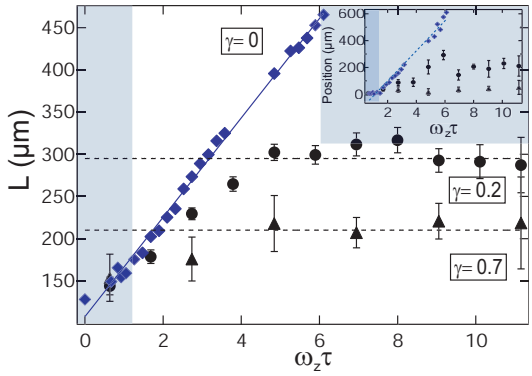

Figure 1: Time evolution of the axial rms size L of the BEC, for various amplitudes σ_V of the random potential, all smaller than the chemical potential μ [$\gamma = \sigma_V/\mu < 1$: γ=0 ($\diamond$), 0.2 ($\bullet$) and 0.7 ($\blacktriangle$)]. The axial trapping frequency is initially $\omega_z/2\pi = 6.7$ Hz and is relaxed during the first 30 ms of the expansion time (grey band). Each point corresponds to an average over three measurements; error bars represent one standard deviation. Inset: Motion of the center of mass of the BEC during the axial expansion for the same values of σ_V. Both sets of data show a strong suppression of transport of the BEC in the presence of disorder.

Fig. 1). The final rms size decreases as the standard deviation σ_V of the random potential increases. The same effect has been observed for various realizations of the random potential. We also observe that the center of mass motion provoked by a longitudinal magnetic 'kick' at the time of release is strongly damped and is stopped in about the same time (see Fig. 1). These observations are *not* made in a regime of tight binding, *i.e.* we observe this localization effect [31] for amplitudes of the random potential small compared to the chemical potential. One may wonder whether our observations can be interpreted in terms of Anderson localization [2]. In fact, in our situation, the interaction energy plays a crucial role, and the healing length is smaller than the typical distance between the speckle grains. This implies a different scenario, which we discuss in this Letter.

We generate an elongated ^{87}Rb BEC in an iron-core electromagnet Ioffe-Pritchard trap [22, 23] with oscillation frequencies, $\omega_\perp/2\pi = 660(4)\,\mathrm{Hz}$ radially and $\omega_z/2\pi = 6.70(7)\,\mathrm{Hz}$ axially. BECs of typically 3.5×10^5 atoms are obtained, with Thomas-Fermi half-length $L_{\mathrm{TF}} = 150\,\mu\mathrm{m}$ and radius $R_{\mathrm{TF}} = 1.5\,\mu\mathrm{m}$, and chemical potential $\mu/2\pi\hbar \sim 5\,\mathrm{kHz}$ [32]. The random potential is turned on at the end of the evaporation ramp and we wait a further 200 ms in the presence of the rf field to ensure that the BEC is in equilibrium in the combined harmonic plus random potential at the end of the sequence.

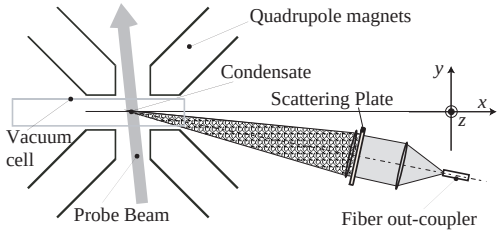

Figure 2: Optical setup used to create the random speckle potential. The long axis of the condensate is oriented along the z direction.

To create the random potential, a $P \leq 150\,\mathrm{mW}$ blue detuned laser beam with optical wavelength $\lambda \simeq 780\,\mathrm{nm}$ is shone onto the atoms, perpendicular to the long axis of the BEC, through a scattering plate (Fig. 2). Positioning the scattering plate as shown in Fig. 2 projects a speckle pattern [26] on the magnetic trap. The scattered beam diverges to an rms radius of 1.83 mm at the BEC.

A speckle pattern is defined by [26] (i) a random intensity $I(\mathbf{r})$ with exponential statistical distribution for which the standard deviation equals the average intensity $\sigma_I = \langle I \rangle$ and (ii) an intensity correlation length Δz, defined as the 'half-width' of the autocorrelation function [26]:

$$\Delta z = 1.22\ \lambda l/D, \tag{1}$$

where D is the beam diameter at the scattering plate and l is the focal length of the lens. We observe the intensity distribution of our speckle pattern on a CCD camera placed at the same distance as the atoms. Taking images with different beam diameters D, we determine the autocorrelation function to obtain the grain size Δz. Taking into account the modulation transfer function (MTF) [27] of the camera, we find that the measured grain size follows Eq. (1). For our setup ($l = 140(5)\,\mathrm{mm}$ and $D = 25.4(1)\,\mathrm{mm}$), we find that $\Delta z = 5.2(2)\ \mu\mathrm{m}$. This is an order of magnitude greater than the healing length $\xi = (8\pi n a)^{-1/2} = 0.11\ \mu\mathrm{m}$ of the trapped condensate. Since $R_{\mathrm{TF}} < \Delta z \ll L_{\mathrm{TF}}$, the optical potential is effectively 1D, with the trapped condensate spread over about $45 - 50$ wells in the longitudinal direction. We characterize the amplitude of the random potential σ_V with respect to the chemical potential μ by [28]:

$$\gamma = \frac{\sigma_V}{\mu} = \frac{2}{3}\frac{\Gamma^2}{2\delta}\frac{\sigma_I}{I_{\mathrm{S}}}\frac{1}{\mu} = \frac{1}{\overline{\omega}}\frac{\Gamma^2}{6\delta}\frac{\sigma_I}{I_{\mathrm{S}}}\left(\frac{15aN}{\overline{a_{\mathrm{ho}}}}\right)^{-2/5} \tag{2}$$

with $\overline{\omega}=(\omega_z\omega_\perp^2)^{1/3}$ and $\overline{a_{\mathrm{ho}}}=(\hbar/m\overline{\omega})^{1/2}$, m the atomic mass, N the condensate atom number, $I_{\mathrm{S}} = 16.56\,\mathrm{W/m^2}$ the saturation intensity, $\Gamma/2\pi = 6.01\,\mathrm{MHz}$ the linewidth, and $a = 5.31\,\mathrm{nm}$ the scattering length. The laser detuning δ we use varies between 0.15 nm and 0.39 nm. The factor 2/3 accounts for the transition strength for π-polarized light. Taking into account our calibration uncertainty, we estimate γ within $\pm 20\%$.

To study the coherent transport of the condensate wavefunction in the random potential, we open the magnetic trap in the axial direction while keeping the transverse confinement and the random potential unchanged. After lowering the current in the axial excitation coils, the axial trapping frequency $\omega_z/2\pi$ is significantly smaller than 1 Hz [33]. The trap is opened in 30 ms, in order to minimize trap loss and heating. Once the current in the axial coils has reached its final value we have a condensate of $N \sim 2.5 \times 10^5 - 3 \times 10^5$ atoms in the magnetic guide.

Our measurements of the axial expansion are shown in Fig. 1. After a total expansion time $t = \tau$ (which includes the 30 ms opening time), we turn off all remaining fields (including the random potential) and wait a further 15 ms of free fall before imaging the atoms by absorption. During this time-of-flight, the axial rms size of the condensate does not increase more than 5%. From profiles of the absorption images we evaluate the axial rms size L of the expanding condensate. In Fig. 1, we plot L versus the axial expansion time τ. In the absence of the random potential ($\gamma = 0$), we observe that the rms size of the atomic cloud grows linearly with $v_{\mathrm{RMS}} \sim 2.47(3)\,\mathrm{mm\,s^{-1}}$ in agreement with the scaling theory [29]. In the presence of the random potential, the expansion dynamics changes dramatically. For a sufficiently high amplitude, the expansion of the cloud is significantly reduced and the

condensate eventually stops expanding. In addition, we observe the damping of longitudinal motion of the center of mass of the cloud (see inset of Fig. 1). This motion is triggered by an axial magnetic 'kick' during the opening of the trap.

These results show a transition from non-inhibited to inhibited transport as the amplitude of the speckle intensity is increased. We studied this phenomenon in further detail by measuring the rms size of the condensate after a fixed axial expansion time ($\omega_z\tau = 4.84$) for different amplitudes σ_V of the random potential. The results are shown in Fig. 3. We see that above a value $\gamma = 0.15$, the rms size decreases with σ_V.

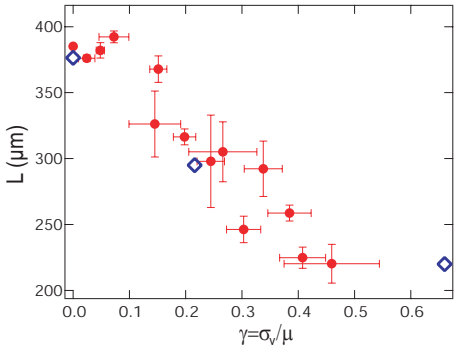

Figure 3: Rms size L of the condensate versus γ after an axial expansion time $\omega_z\tau = 4.84$. The points $\diamond$ correspond to the curves of Fig. 1.

From the absorption images, we also evaluate the density in the magnetic guide, after correcting for radial expansion during the time-of-flight. We observe that the density at the center of the BEC does not drop by more than a factor of 2 for $\gamma > 0.2$. Therefore, we conclude that the interaction energy dominates at the centre of the condensate trapped by disorder, a point we shall discuss below.

To understand the suppression of expansion of the BEC in the random potential, we have performed numerical calculations of the BEC dynamics in the mean-field Gross-Pitaevskii approach. We consider a BEC trapped in a cylindrically-symmetric 3D-harmonic trap with frequencies $\omega_\perp$ and ω_z in the radial and axial directions respectively. Assuming tight confinement in the radial directions ($\hbar\omega_\perp \gg \hbar\omega_z, \mu, k_BT$), the dynamics is reduced to one-dimension. In addition, the BEC is subjected to a static random potential $V(z)$ which corresponds to a numerically-generated speckle field [26] $V(z) = \sigma_V v(z)$ where $v(z)$ is a normalized speckle pattern with $\langle v\rangle^2 = \langle v^2\rangle/2 = 1$. We consider parameters close to the experimental situation described above. In particular, the healing length ($\xi \simeq 8\times10^{-4}L_{\rm TF}$) and the speckle correlation length ($\Delta z \simeq 0.049 L_{\rm TF}$) are much smaller than the size of the BEC.

We first compute the static 1D BEC wavefunction in the combined (harmonic plus random) trap. Because $\xi \ll \Delta z$, the density profile simply follows the modulations of the combined trap in the Thomas-Fermi regime $|\psi(z)|^2 = [\mu - m\omega_z^2 z^2/2 - V(z)]/g_{1D}$ in the region where $\mu > m\omega_z^2 z^2/2 + V(z)$ and $|\psi(z)|^2 = 0$ elsewhere. Here, m is the atomic mass and $g_{1D} = 2\hbar a\omega_\perp$ the 1D interaction parameter. At time $t = 0$, we suddenly switch off the harmonic confinement in the axial direction while keeping unchanged the interaction parameter g_{1D} and the random potential and we compute the time-evolution of the BEC. The results for the axial rms size L of the BEC are plotted in Fig. 4a for various amplitudes of the random potential $\sigma_V = \gamma\mu$. In the absence of disorder, the evolution of the BEC corresponds to self-similar expansion with parameter $b(t) \sim \sqrt{2}\omega_z t$ [34]. For small amplitudes of disorder ($\gamma \lesssim 0.1$), the BEC grows by about one order of magnitude for $\omega_z t = 10$. For larger amplitudes of disorder ($\gamma \gtrsim 0.15$), after initial expansion, the BEC stops expanding [35]. This is consistent with the experimental findings. This strong suppression of expansion corresponds to disorder-induced trapping of the BEC.

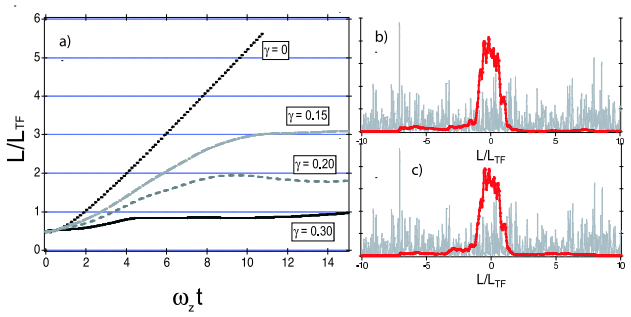

Figure 4: a) Time evolution of the rms size L of the BEC in the random potential V for various values of the speckle amplitude $\sigma_V = \gamma\mu$ and comparison to the theoretical prediction in the absence of random potential (black dotted line). b-c) Density profiles of the BEC trapped by disorder at two different times $\omega_z t = 10$ and $\omega_z t = 20$ respectively.

We now describe a scenario for disorder-induced trapping of the BEC when released from the harmonic trap. For small enough random potential, the initial stage of expansion can be described using the scaling theory [29]. According to this theory, the fast atoms populate the wings of the expanding BEC whereas the slow atoms are close to the center. It is thus tempting to distinguish two regions of the BEC: (i) the central part where the interaction energy dominates the kinetic energy and trapping is due to the competition between interactions and disorder, and (ii) the wings where the kinetic energy exceeds the interactions and trapping is rather due to the competition between the kinetic energy and disorder.

In the center, the average density and thus the effective chemical potential $\overline{\mu}$ slowly decrease during the expan-

sion stage. As the interaction energy is much larger than the kinetic energy, the local density adiabatically follows the instantaneous value of $\overline{\mu}$ in the Thomas-Fermi regime $|\psi(z)|^2 = [\overline{\mu} - V(z)]/g_{1D}$ in the region where $\overline{\mu} > V(z)$ and $|\psi(z)|^2 = 0$ elsewhere. This evolution stops with fragmentation, *i.e.* when the BEC meets two peaks of the random potential with amplitudes larger than $\overline{\mu}$. Using the statistical properties of the random potential [26], we can estimate the probability of such large peaks and we conclude that this happens when the central density n_0 reaches the value

$$n_0 \simeq 1.25 \left(\frac{\sigma_V}{g_{1D}}\right) \ln\left(\frac{0.47 L_{TF}}{\Delta z}\right) . \qquad (3)$$

This formula is in good agreement with our numerical and experimental findings [30].

Due to the small density, the situation is completely different in the wings which are populated by almost free particles interacting with the disordered potential. The BEC wavefunction thus undergoes disorder-induced multiple reflections and transmissions and is ultimately blocked by a large peak of the speckle potential. Therefore, the BEC is not in the Thomas-Fermi regime and the local density is not stationary. Due to the conservation of the energy, the kinetic energy per particle ϵ is of the order of the typical energy in the initial condensate ($\epsilon \sim \mu$) so that the typical wavelength of the fluctuations in the wings is of the order of the healing length in the initial condensate $\lambda_w \sim \xi = \hbar/\sqrt{2m\mu}$.

This scenario of disorder-induced trapping is accurately supported by our numerical integration of the Gross-Pitaevskii equation. In particular, the density profiles plotted in Fig. 4 show the static Thomas-Fermi shape in the center and time-dependent fluctuations in the wings with typical wavelength $\lambda_w \sim \xi$ [30].

In conclusion, we have experimentally investigated transport properties of an interacting BEC in a random potential. Controlling the strength of disorder, we have observed the transition from free expansion to absence of diffusion as disorder increases. We have presented numerical simulations that reproduce well the observed suppression of expansion and we have discussed a theoretical model that describes the scenario for disorder-induced trapping. In the future, it would be interesting to further investigate this highly controllable system, for example by changing the correlation length of disorder or employing Bragg spectroscopy to probe the momentum spectrum of the BEC [23].

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[33] From the variations of the dipole curvature with the current, we estimate the upper bound of $\omega_z/2\pi$ to be ~ 400 mHz, which is compatible with the linear expansion observed in the absence of disorder (see Fig. 1).

[34] Our findings are in excellent agreement with the analytic solution of the scaling equations of [29], $\sqrt{b(t)(b(t)-1)} + \ln[\sqrt{b(t)} + \sqrt{b(t)-1}] = \sqrt{2}\omega_z t$.

[35] It should be noted that for $\gamma \gtrsim 0.25$, the fluctuations of Δz are due to small still evolving contributions in the wings of the BEC wavefunction while the core of the wavefunction is localized (see below).