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Tunable source of correlated atom beams

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We use a one-dimensional optical lattice to modify the dispersion relation of atomic matter waves. Four-wave mixing in this situation produces atom pairs in two well defined beams. We show that these beams present a narrow momentum correlation, that their momenta are precisely tunable, and that this pair source can be operated in the regimes of low mode occupancy and of high mode occupancy.

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In quantum optics, existence of mechanisms to produce photon pairs, such as parametric down-conversion, enabled the realization of several fundamental experiments on quantum mechanics. For example, the violation of Bell’s inequalities [1] or the Hong-Ou-Mandel effect [2] reveal the surprising properties of quantum correlations in entangled photon pairs. These fascinating properties have found applications in quantum information and communications [3]. In analogy to photon pairs, there have been several recent demonstrations of correlated atom pairs production [4–10]. In particular, momentum correlations of spatially separated samples is an important requirement for the demonstration of an atomic Einstein-Podolsky-Rosen state [11, 12] and the violation of Bell’s inequalities. Such momentum correlations were demonstrated for atom pairs produced by molecule dissociation [4] or by spontaneous four-wave mixing in free space through the collision of two Bose-Einstein condensates (BECs) [5, 13]. In these experiments the pairs which were produced lay on a spherical shell. This geometry is disadvantageous because many spatial modes are populated, and if one wishes to use Bragg diffraction to manipulate and recombine the pairs on a beam splitter [11, 14], the vast majority of the pairs is unusable.

On the other hand, if pair production is concentrated in a small number of modes, experimenters can make more efficient use of the generated pairs. One can then choose to work either with low mode occupation, the well separated pair regime, or with high mode occupation, referred to as the squeezing regime in Ref. [15]. An example of twin beams generated in the latter regime is described in Ref. [6]. The squeezing regime is well suited to the study of highly entangled multiparticle systems and for investigations of atom interferometry below the standard quantum limit [16, 17]. The source we study in this Rapid Communication can be operated in both regimes. We use atomic four-wave mixing in a one-dimensional (1D) optical lattice, which results in production of atom pairs

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We thus control the value of $k_0/k_{\text{rec}} = h \delta \nu / 4E_{\text{rec}}$, the BEC’s momentum in the lattice frame. The lattice is held on for a duration $T_L = 2$ ms, and suddenly switched off, simultaneously with the optical trap. To avoid magnetic perturbation of the cloud during freefall, we apply an RF pulse that transfers 50% of the atoms to the field insensitive $m_x = 0$ sublevel [26]. The atoms remaining in $m_x = 1$ are subsequently removed by a strong magnetic gradient. After a 307 ms mean time of flight, the $m_x = 0$ atoms fall on a microchannel plate detector, which permits 3D reconstruction of the atomic cloud [28].

As shown in Fig. 1(b), we observe three main density peaks after the time of flight. The tallest is the initial BEC. The two others are formed by atoms scattered into momentum classes centered in $k_1$ and $k_2$, whose values are consistent with those expected from the phase-matching conditions illustrated in Fig. 1(a). Since the optical lattice is switched off abruptly, the Bloch states of momenta $k_0$, $k_1$ and $k_2$ are projected onto plane waves, mainly in the first Brillouin zone due to the low lattice depth. Each of the beams at $k_1$ and $k_2$ contains about $10^3$ detected atoms, which we estimate to correspond to about $2 \times 10^3$ atoms per beam. We also detect some atoms between the beams, which result from scattering into excited transverse modes [29]. Due to the low overlap between the transversely excited states and the initial wave function, this transverse excitation is far less efficient than the previously described 1D process. In addition, scattered atoms can also undergo secondary scattering contributing to the background between the beams.

In the following, we focus on the two beams. Using them for quantum atom optics experiments or for interferometry will require recombinging them. It is therefore crucial to know the width of their correlation. From the 3D-momentum distribution $n(k)$, we computed the normalized second-order cross correlation function,

$$g_C^{(2)}(k, k') = \frac{\langle n(k)n(k') \rangle}{\langle n(k) \rangle \langle n(k') \rangle}$$

where $k$ belongs to beam 1 and $k'$ to beam 2. The BEC is not exactly at rest in the optical trap, but exhibits shot-to-shot momentum fluctuations on the order of $10^{-2} k_{\text{rec}}$. We correct for these fluctuations by recentering separately the single shot momentum distributions $n(k)$ around $k_1$ and $k_2$, using the shift obtained from Gaussian fits to the peak at $k_1$ and to the diffraction peak at $k_0 + 2k_{\text{rec}}$. This correlation function exhibits a peak for $k_z \approx k_1$ and $k'_z \approx k_2$ [Figs. 3(a) and 3(b)]. The presence of this peak indicates that the two atomic beams are indeed correlated.

We wish to determine the number of modes present in each beam, and how many of these modes are correlated. We therefore examine the local second-order correlation function of a single beam, $g_L^{(2)}(k, k')$, which is obtained as in Eq. (1) but with both $k$ and $k'$ belonging to beam 1. This correlation function, plotted in Figs. 3(c) and 3(d), exhibits bunching for $k'_z \approx k_z \approx k_1$, due to den-
FIG. 3. (a,b) Cuts along \((y,z)\) of the integrated, normalized cross correlation function of the two beams, \(g^{(2)}(\Delta k) = f dk_1 g^{(2)}(k_1, k_2 + \Delta k)\). The integration over the momentum distribution \(k_1\) is performed on a box with dimensions \(L_{k_x} = L_{k_y} = 0.4 k_{\text{rec}}\) and \(L_{k_z} = 5 \times 10^{-2} k_{\text{rec}}\) centered on beam 1, \(k_1 + k_2 = (k_1 + k_2) k_{\text{rec}}\) and the cuts have a thickness \((1.5 \times 10^{-1} k_{\text{rec}})\) along \(z\) \((x\) and \(y)\). The bunching, due to the correlation between the two beams, has a longitudinal (transverse) width \(\sigma_{\text{e},z} = 1.8 \times 10^{-2} k_{\text{rec}}\) \((\sigma_{\text{e},y} = 1.6 \times 10^{-1} k_{\text{rec}})\).

(c,d) Cuts along \((y,z)\) of the integrated, normalized local correlation function of beam 1, \(g^{(2)}(\Delta k) = f dk_1 g^{(2)}(k_1, k_1 + \Delta k)\). The integration region is the same as for the cross correlation, and the cuts have a thickness \(2.5 \times 10^{-3} k_{\text{rec}}\) \((0.1 k_{\text{rec}})\) along \(z\) \((x\) and \(y)\). The bunching, due to HBT effect, has a longitudinal (transverse) width \(\sigma_{\text{e},z} = 3.7 \times 10^{-3} k_{\text{rec}}\) \((\sigma_{\text{e},y} = 1.3 \times 10^{-1} k_{\text{rec}})\). Cuts along \(x\) (not shown here) have same widths and amplitudes as cuts along \(y\). These correlation functions are calculated using 850 experimental realisations, with \(k_{\text{rec}} = 0.65 k_{\text{rec}}\), a lattice depth \(V_0 = 0.725 E_{\text{rec}}\) and a lattice duration \(T_L = 2\) ms. In all plots, the horizontal error bars indicate the bin size and the vertical ones correspond to the statistical \(1 \sigma\) uncertainties. The solid lines are Gaussian fits to the data from which we extract the correlation widths.

It appears in Fig. 3 that, while in the transverse direction, the cross and local correlations have similar widths [Figs. 3(a) and 3(c)], the cross correlation is 5 times broader than the local one along the vertical axis [Figs. 3(b) and 3(d)]; each mode is correlated with several modes of the other beam. If one uses two such beams as inputs to a beam splitter, this broadening amounts to a loss of coherence, and the interference contrast would be reduced. We emphasize that the observed widths may be broadened by other effects, and so their numerical ratio is not exactly equal to the number of correlated modes. For the local correlation, we estimate that the finite vertical resolution of the microchannel plate detector contributes notably to the observed width. This resolution comes about because the surface which defines the atom arrival time is not flat but consists of tilted channels which intercept the atoms at different heights. The width shown in Fig. 3(d) is consistent with this interpretation. For the cross correlation, the observed width is broadened by the fact that the vertical source size is not negligible [32]. Note also that the limited coherence of the initial quasi-BEC plays a role in the cross correlation width [32].

The use of an optical lattice permits control over the output beam momenta. Changing the detuning \(\delta\nu\) between the lattice beams results in varying the value of \(k_0\). In Fig. 4, we plot the mean vertical momenta \(k_1\) and \(k_2\) of both beams, measured for different \(k_0\), as well as the expectation (solid line) based on the phase-matching conditions illustrated in Fig. 1(a). We obtain a fair agreement over a large range, even though the solid line presents a small shift in comparison to the data points and does not reproduce the observed shape for high values of \(k_0\). However, as already observed for four-wave mixing in free space [33], phase-matching conditions can be influenced by mean-field effects. A simple correction to the phase-matching curve is found just by adding the mean field to the energy conservation condition: Since the two atoms of a scattered pair are distinguishable from the atoms...
of the initial BEC, the mean-field energy experienced by each of them is not \( g n_0 \) (with \( g = 4\pi\hbar^2 a / m \), \( a \) and \( m \) the scattering length and the mass of He* and \( n_0 = 10^{13} \) atoms/cm\(^3\) the BEC’s density), but \( 2g n_0 \), so that the energy conservation condition reads:

\[
2E(k_0) + 2gn_0 = E(k_1) + E(k_2) + 4gn_0 \tag{2}
\]

where the energy \( E(k) \) is given by the dispersion relation in the first Bloch band of the lattice without interaction. As seen in Fig. 4 (dashed line), this correction leads to very good agreement with the experimental data, and accounts for the shift of the phase-matching curve and the change of its shape. A more exact calculation of the phase-matching conditions, inspired by Ref. [21], confirms the accuracy of Eq. (2) in our experimental conditions and will be given in Ref. [34].

Another degree of freedom results from the fact that pair creation only takes place while the lattice is on. We can thus tune the beam populations with the lattice duration \( T_L \). In the example of Fig. 5 these populations increase exponentially with \( T_L \) during a few hundreds of \( \mu s \), and then reach a plateau. This saturation could be explained by several mechanisms such as the decrease of spatial overlap between condensate and scattered beams [19], multimode effects [35] and secondary scatterings from the beams. Condensate depletion is at most about 20 %, and should be of little importance in the saturation. For small \( T_L \), there is no discernable population difference between both beams. By contrast, we observe that at large \( T_L \) the population of beam 1 is almost twice that of beam 2, a phenomenon also noticed in Ref. [19]. This may be due to \( k_2 \) being in a dynamically unstable region while atoms with quasimomentum \( k_1 \) can only undergo secondary scattering to excited transverse modes.

At intermediate \( T_L \), we observe negligible losses due to secondary scattering and high mode population (around 60 atoms per mode at \( T_L = 0.2 \) ms in the example of Fig. 5). The resulting beams should contain strongly correlated pairs. In an attempt to verify a nonclassical correlation, we examined atom number difference between the two beams. By selecting two regions around the centers of the two beams, we do indeed observe a sub-Poissonian number difference [6, 31], as shown in Fig. 6. The observed variance is consistent with that observed in Ref. [31], and is limited in large part by the quantum efficiency of the detector. Other features of the variance are puzzling, however. First the minimum of the dip in the variance occurs when the center of region 1 is shifted by 0.1 \( k_{rec} \) with respect to the center of the density distribution in beam 1. Second, in the transverse plane, the size of the regions over which the variance is reduced is nearly an order of magnitude smaller than the transverse width of the correlation function. We plan to investigate these effects in future experiments.

To conclude, we have demonstrated an efficient process for the production of correlated atom pairs. We have control over both the final momenta and the intensity of the correlated beams. We characterize the width of the correlation in momentum and find evidence of sub-Poissonian fluctuations of population difference. This source should be useful in multiple particle interference experiments both in the regime of well-isolated pairs [12] and in the regime of large occupation numbers [11].

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[14] In all this paper we denote momenta by their value divided by \( \hbar \).