

Twin-atom beams

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In recent years, substantial progress has been made in exploring and exploiting the analogy between classical light and matter waves for fundamental investigations and applications¹. Extending this analogy to quantum matter-wave optics is promoted by the nonlinearities intrinsic to interacting particles and is a stepping stone towards non-classical states^{2,3}. In light optics, twin-photon beams⁴ are a key element for generating the non-local correlations and entanglement required for applications such as precision metrology and quantum communication⁵. Similar sources for massive particles have so far been limited by the multi-mode character of the processes involved or a predominant background signal^{6–13}. Here we present highly efficient emission of twin-atom beams into a single transversal mode of a waveguide potential. The source is a one-dimensional degenerate Bose gas¹⁴ in the first radially excited state. We directly measure a suppression of fluctuations in the atom number difference between the beams to 0.37(3) with respect to the classical expectation, equivalent to 0.11(2) after correcting for detection noise. Our results underline the potential of ultracold atomic gases as sources for quantum matter-wave optics and should enable the implementation of schemes previously unattainable with massive particles^{5,15–19}.

Binary collisions between atoms provide a natural means to generate dual number states of intrinsically correlated atoms¹⁵. Experimental schemes include spontaneous emission of atom pairs by collisional de-excitation⁹ or four-wave mixing^{10–12}. Stimulated emission into twin-modes has been demonstrated in seeded four-wave mixing^{2,6}, and parametric amplification in optical lattices^{7,8} or spinor condensates¹³. Among these schemes, suppression of relative number fluctuations could so far be demonstrated only for multi-mode twin-atoms¹². A different route to non-classical states is provided by ensembles in multi-well potentials that become number-squeezed during their time evolution^{20,21}.

Here, we demonstrate how collisional de-excitation of a one-dimensional degenerate Bose gas can be used to efficiently create matter-wave beams of twin-atoms. The restricted geometry of a waveguide potential forces emission of the beams into a single transversal mode, in analogy to an optical parametric amplifier⁴. We prepare the initial population inversion to a radially excited state by shaking the trap, employing an optimal control strategy. Time-of-flight fluorescence imaging is used to directly observe the suppressed relative number fluctuations in the emitted beams.

The starting point of our investigations is a dilute, quantum degenerate gas of neutral rubidium-87 atoms magnetically trapped in a tight waveguide potential with a shallow axial harmonic confinement ($\nu_x = 16.3$ Hz) on an atom chip²². Our scheme

relies on an effective two-level system in the radial vibrational eigenstates of the waveguide. This is accomplished by creating unequal level spacings in the radial yz -plane by radio frequency dressing²³, which introduces anharmonicity and anisotropy. The resulting single-particle first and second excited state energies are $E_{y,z}^{(1)} = h \cdot [1.83, 2.58]$ kHz and $E_{y,z}^{(2)} = h \cdot [3.82, 5.22]$ kHz. As the level spacings increase at higher levels, the ground state $|n_y, n_z; k_x\rangle = |0, 0; 0\rangle$ ($n_{y,z}$ and k_x denoting the radial quantum numbers and the axial momentum, respectively) and the first excited state along y , $|1, 0; 0\rangle$ have the lowest energy difference among all possible combinations (Fig. 1a), establishing a closed two-level system.

Using standard techniques we generate a Bose gas of typically 700 atoms at a temperature $T \lesssim 40$ nK $\approx h/k_B \cdot 830$ Hz (obtained independently from fits to a non-excited degenerate gas²⁴ and its residual thermal fraction²⁵). The thermal occupation of state $|1, 0; 0\rangle$ is negligible, and the chemical potential quantifying the mean-field interaction is $\mu \sim h \cdot 500$ Hz $\ll E_y^{(1)}$. Our Rb sample is therefore a one-dimensional, weakly interacting quasi-Bose-Einstein condensate (quasi-BEC, Lieb-Liniger $\gamma \sim 0.008$; ref. 14), which can be seen as a locally coherent matter wave²⁶ with a coherence length approximately one order of magnitude below the system size.

Having prepared the gas, we create a population inversion by transferring the quasi-BEC almost entirely to state $|1, 0; 0\rangle$ (see Fig. 1a). The transition is driven by shaking the trap along the radial y -direction on the scale of the ground state size (~ 100 nm). The trajectory (total duration 5 ms, see Fig. 2a) has been optimized by employing an iterative optimal control algorithm (see Supplementary Information). In the experiment, the displacement is achieved by driving a current in an auxiliary chip wire parallel to the main trapping wire.

We monitor the radial momentum distribution of the quasi-BEC by releasing the cloud from the trapping potential at different times t , during and after the excitation pulse. Images are taken after 46 ms of ballistic expansion (Fig. 2b), using a single-atom-sensitive fluorescence imaging system²⁷. After the excitation, we observe a small residual beating between the macroscopically occupied $|1, 0; 0\rangle$ state and a remaining non-excited population in $|0, 0; 0\rangle$. From a fit to the beating pattern (Fig. 2c) we estimate an efficiency of the coherent transfer of $\eta_e \approx 97\%$ and deduce the energy difference $\epsilon = h \cdot 1.78$ kHz between $|0, 0; 0\rangle$ and $|1, 0; 0\rangle$ (see Methods). The slight deviation between $E_y^{(1)}$ and ϵ is explained by particle interactions. A calculation based on the one-dimensional Gross-Pitaevskii equation (GPE) (Fig. 2d) shows excellent agreement with the observed dynamics.

The population inversion to $|1, 0; 0\rangle$ represents a highly non-equilibrium state of the system, analogous to a laser gain medium

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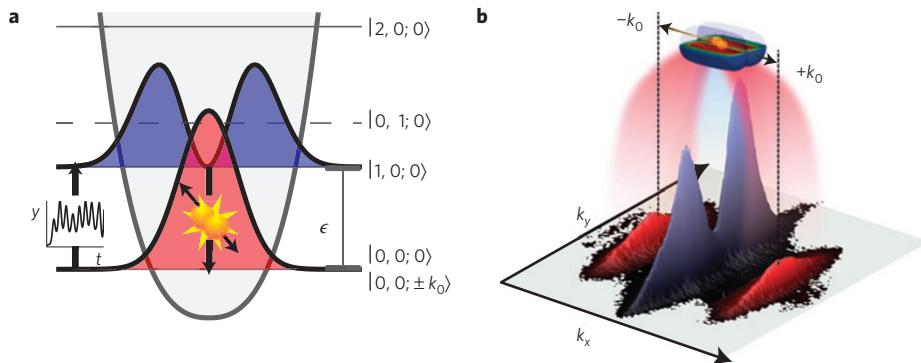


Figure 1 | Schematic of the excitation and emission process. **a**, The quasi-BEC is transferred from the ground state $|0, 0; 0\rangle$ into $|1, 0; 0\rangle$, the first excited state of the trapping potential along the radial y -direction. This is accomplished by means of a fast non-adiabatic movement of the potential minimum along an optimized trajectory (inset). The excited state decays by emission of twin-atoms into the radial ground state modes $|0, 0; \pm k_0\rangle$. **b**, After excitation and pair emission, the cloud is released from the trapping potential and imaged after expansion. The central part of the system clearly shows the spatial structure of the radially excited state (blue). Two clouds containing the twin-atoms (red) are emitted.

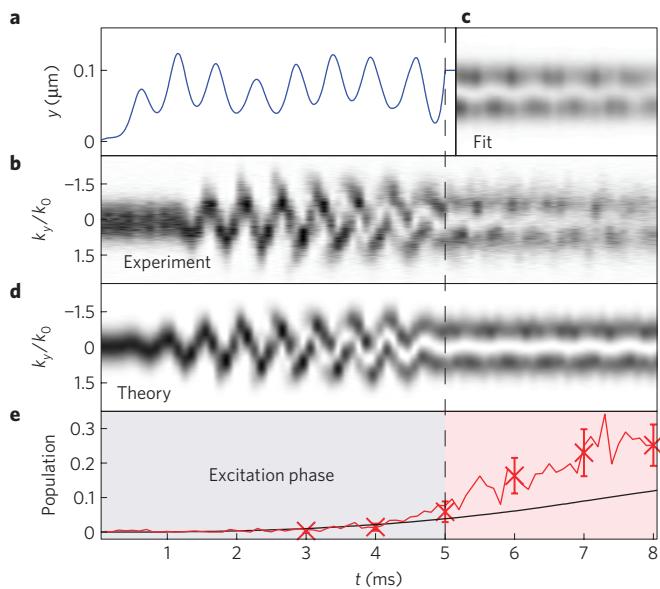


Figure 2 | Dynamics of the excitation and emission process: comparison between theory and experiment. **a**, Optimized control trajectory for the trap movement along y . **b**, Measured momentum distribution. Fluorescence images integrated along k_x over a region encompassing the central cloud. Average over five experimental runs. **c**, Fit of the time-dependent momentum distribution, scaled to the experimental total density at each time step. **d**, Calculated momentum distribution, using a one-dimensional GPE model. As the model does not take into account the pair emission process, agreement with the experiment is expected approximately up to the end of the excitation pulse. **e**, Red line: Population of the emitted clouds obtained from the same data set as **b**. Red crosses: population of the emitted clouds obtained from a separate measurement with 100 experimental runs each (error bars represent the ensemble standard deviation). Black line: theoretical estimation for spontaneous processes only.

after a pump pulse. For the ensuing relaxation, the only allowed channel is a two-particle collisional process, emitting atom pairs with opposite momenta. In contrast to experiments with free-space collisions^{10,12} or two-dimensional gases⁹, the constricted geometry and non-degenerate level scheme of our source restricts the outgoing matter waves to the radial ground state of the waveguide, yielding twin-atom beams in a single transversal mode. Within a binary collision, two atoms are scattered from $|1, 0; 0\rangle|1, 0; 0\rangle$

to $|0, 0; +k_0\rangle|0, 0; -k_0\rangle$, where energy conservation requires the final momenta to be centred around $\pm k_0 = \pm\sqrt{2m\epsilon/\hbar}$. The emission process can be understood as a matter-wave analogue of a degenerate optical parametric amplifier, where the initially empty twin-modes are seeded by vacuum fluctuations⁴ and have an exponentially growing population if phase matching conditions are fulfilled. Owing to the finite size, the axial multi-mode character of the quasi-BEC source¹⁴, and the depletion of the initial state, a comprehensive understanding of the process is challenging. The emission dynamics, as well as the ensuing longitudinal coherence properties of the beams (which should resemble those of the source), will be the subject of future experimental and theoretical studies. A comparison of the observed rates to a simple calculation using Fermi's golden rule (see Supplementary Information) is shown in Fig. 2e and demonstrates the insufficiency of a purely spontaneous model (in contrast to the findings of ref. 9 for a transversal multi-mode system).

Once the trap potential is switched off and the atoms propagate freely (Fig. 1b), the twin-beam modes can be detected essentially background-free in fluorescence images (Fig. 3a,b) as they separate from the source. In Fig. 3c, the radial momentum distribution of the twin-beams is compared to an independent measurement of the initial $|0, 0; 0\rangle$ cloud. The small deviation is attributed to a slight overlap of the central cloud with the integration regions of the emitted clouds (red boxes in Fig. 3b) and interactions with the mean field of the quasi-BEC in the $|1, 0; 0\rangle$ state. Furthermore, an excited cloud at $t = 6$ ms is shown. From the axial position of the side peaks (Fig. 3d) we can deduce an emitted atom momentum of $k_0 = 2\pi \cdot 0.883(3) \mu\text{m}^{-1}$, equivalent to $\epsilon = h \cdot 1.78(1) \text{ kHz}$, in perfect agreement with the value determined from the beating fit (Fig. 2c). The width of the emitted clouds (Fig. 3d) is increased by a factor of ≈ 1.4 with respect to the source cloud. At momenta corresponding to $\epsilon' \approx h \cdot 3.9 \text{ kHz}$, further, very weakly populated atom clouds ($\lesssim 1$ atom per image) are observed on an averaged picture (Fig. 3b). They imply a transient population of $|2, 0; 0\rangle$ during the control sequence that directly decays into the radial ground state. In the following analysis, they are merged with the atoms originating from $|1, 0; 0\rangle$.

The non-classical correlation in the emitted twin-atom beams is revealed by a sub-binomial distribution of the number imbalance $n = N_1 - N_2$ between atoms detected at $\pm k_0$. The variance of n can be expressed as $\sigma_n^2 = \xi^2 \bar{N}$, where \bar{N} denotes the mean total atom number in the emitted clouds. The noise reduction factor ξ^2 quantifies the suppression of σ_n^2 with respect to a binomial distribution and, thus, the amount of correlation between the populations N_1 and N_2 .

In the fluorescence images, we count photons in regions encompassing the emitted clouds that were released at $t = 7$ ms.

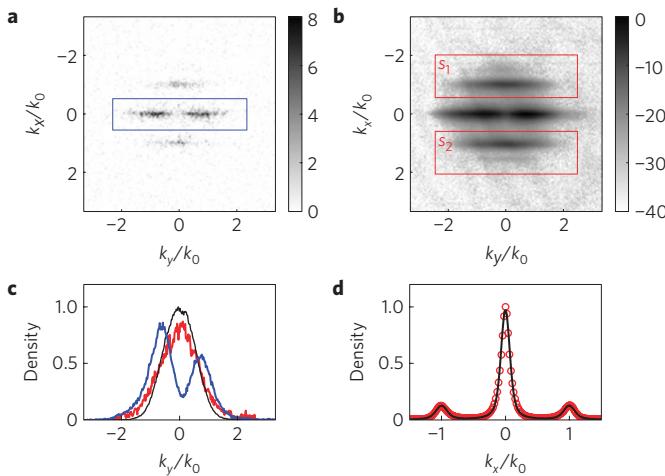


Figure 3 | Atom cloud image analysis. **a**, Typical experimental image of ~ 700 atoms released from the trap 7 ms after starting the excitation sequence. The cloud is allowed to expand for 46 ms, making the initial momentum distribution accessible. The quasi-BEC in the excited state $|1, 0; 0\rangle$ is clearly distinct from the emitted clouds at momenta $\pm \hbar k_0$. Units are photons per pixel. The blue box indicates the integration range for the data shown in Fig. 2b. **b**, Average over $\approx 1,500$ images similar to **a**. The colour scale is logarithmic (dB referenced to peak density). The regions used for correlation analysis are indicated as red boxes. **c**, Normalized, radial momentum distributions of the central (blue) and emitted (red) clouds. Average of 50 images of clouds released at $t = 6$ ms. As comparison, the distribution of a non-excited cloud is shown (black, average over 100 images). **d**, Normalized profile of **b** along k_x (red dots) and three-peak fit (black line) based on stochastic simulations²⁴.

For given atom numbers $N_{1,2}$, the expectation values for the photon numbers are $S_{1,2} = \bar{p}N_{1,2} + \bar{b}/2$, where $\bar{p} = 12.3(9)$ denotes the average number of photons per atom and $\bar{b}/2$ accounts for background events. Our main observable is the variance σ_s^2 of the signal imbalance $s = S_1 - S_2$. Its expectation value for a binomial distribution of atoms and noise-free detection is given by $\sigma_{\text{bin}}^2 = \bar{p}\bar{S}$, where $\bar{S} = \bar{S}_1 + \bar{S}_2 - \bar{b}$. From the experimental data as shown in Fig. 4a we obtain an uncorrected reduction factor $\sigma_s^2/\sigma_{\text{bin}}^2 = 0.37(3)$. However, a significant contribution to σ_s^2 originates not from the atom number fluctuations, but from the detection process itself. For fluorescence imaging as employed in our experiment, this contribution can be directly calculated from the photon shot noise and detection background (see Methods). It is accounted for by subtracting a correction σ_d^2 from σ_s^2 . From the corrected variances, we infer a reduction factor of $\xi^2 = (\sigma_s^2 - \sigma_d^2)/\sigma_{\text{bin}}^2 = 0.11(2)$, which is the main result of this paper. It is equivalent to an intensity squeezing in the sense of ref. 4. In strong contrast to the suppressed relative fluctuations, applying an analogous calculation on the variance σ_s^2 of the summed signal in the emitted clouds $S = S_1 + S_2$ (binned into groups with similar total atom number) yields super-Poissonian fluctuations $(\sigma_s^2 - \sigma_d^2)/\bar{p}\bar{S} \sim 7$, highlighting the presence of bosonic amplification.

To study the correlation data in more detail, we bin the experimental shots according to the emitted atom signal S and calculate the variances σ_s^2 and mean signals \bar{S} for each bin (consisting of typically 100 runs) separately, as shown in Fig. 4b. The differences between the data points and the detection noise σ_d^2 represent the corrected variances, as introduced above. They seem to be independent of \bar{S} , which is supported by χ^2 -test results on various plausible models. As any uncorrelated emissions should scale with \bar{S} , this suggests that the non-zero value of ξ^2 can be explained by a slight additional background signal, for

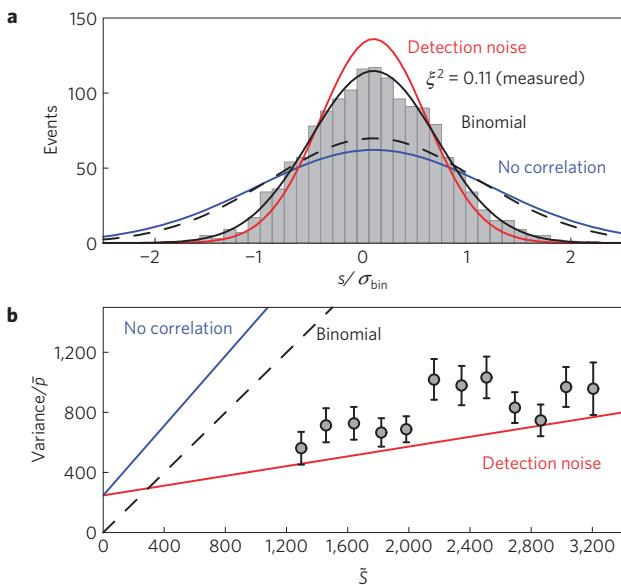


Figure 4 | Correlation analysis. **a**, Histogram of observed signal imbalances s between the emitted clouds, in units of the binomial standard deviation $\sigma_{\text{bin}} = (\bar{p}\bar{S})^{1/2}$. The curves indicate normal distributions corresponding to the experimental result of $\xi^2 = 0.11(2)$ (black, solid); the limits of perfect correlation, where only detection noise remains (red, solid); and a binomial distribution for $p\bar{S}$ trials (black, dashed). **b**, Observed signal imbalance variances for data bins corresponding to different total signal in the emitted clouds \bar{S} . Error bars are the standard error. The lines correspond to those in **a**. The corrected variances are given by the vertical distances between the data points and the detection noise.

example as a result of the residual overlap of the excited quasi-BEC and the emitted clouds.

The availability of single-mode twin-atom beams adds an essential building block for quantum matter-wave optics. As our scheme does not rely on the internal structure of the atoms, it can be applied to any sufficiently controllable system of interacting bosons. Possible applications include interferometry with dual-Fock states¹⁵, Hong–Ou–Mandel-type experiments¹⁷ or continuous-variable entanglement^{5,19}. Inclusion of internal (for example, hyperfine states with appropriate scattering properties) or few-mode external (for example, two-mode double well states) degrees of freedom seems a viable strategy to generate non-local entangled states of massive particles for Bell-type measurements^{16,18,19,28,29}.

Methods

Trap potential preparation. We trap our atoms in an adiabatic dressed state potential²³, created on an atom chip²². The radial potential can be approximated by quartic polynomials, where the x^4 -term typically contributes $\approx 5\%$ at relevant length scales (see Supplementary Information for details).

Radial excitation. We use optimal control of the one-dimensional (along y) GPE (ref. 30) to derive the shape of the excitation pulse. The anharmonicity of the radial potential enables transfer to the population-inverted state, as opposed to a harmonic potential in which only coherent states can be excited by displacement. To obtain the excitation efficiency $\eta_e = 0.97$ and $\epsilon = h \cdot 1.78$ kHz from the beating between $|0, 0; 0\rangle$ and $|1, 0; 0\rangle$ we fit a superposition of the $\psi_1(k_y) = \langle k_y | 1, 0; 0 \rangle$ and $\psi_0(k_y) = \langle k_y | 0, 0; 0 \rangle$ wavefunctions with a time-dependent relative phase $\epsilon/h \cdot t$ to the momentum density of the central cloud (Fig. 2b,c). Starting from about $t \approx 5$ ms, the contrast of the beating pattern starts to fade (Fig. 2b), which is attributed to dephasing collisions between emitted atoms and the main cloud. This is taken into account by (incoherently) adding a contribution $|\psi_0(k_y)|^2$ to the density fit, with a linearly growing population reaching $\approx 28\%$ at the latest time measured. Note, however, that this does not affect the efficiency of the excitation itself.

Imaging characterization. The fluorescence imaging system²⁷ is calibrated by comparing a large number of fluorescence and calibrated absorption images of identically prepared clouds. As absorption images yield an absolute atom number, we can determine the number of photons per atom to $\bar{p} = 12.3(9)$. This value is compatible with independent measurements of the physical properties of three-dimensional clouds²⁵. The detection noise σ_d^2 , which is used to correct the observed number fluctuations, can be split into a constant and a signal-dependent part (see red line in Fig. 4b). The constant part is given by background events and readout noise $\sigma_b^2 \approx 2\bar{f}$, which can be obtained from empty regions in the images next to the atoms. The signal-dependent contribution is defined by the variance of the photon number per atom σ_p^2 . Photon shot noise and amplification noise of the Electron Multiplying Charge Coupled Device (EMCCD) camera used yields $\sigma_p^2 = 2\bar{p}$ (see Supplementary Information and ref. 27). In total, we arrive at $\sigma_d^2 = \sigma_p^2\bar{N} + \sigma_b^2 = 2\bar{S} + \sigma_b^2$.

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Author contributions

R.B., S.M. and T. Berrada collected the data presented in this letter. J.G. and U.H. provided the OCT calculations for the excitation sequence. All authors contributed to analysis and interpretation of the data and helped in editing the manuscript.

Additional information

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