Observation of quantized conductance in neutral matter

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In transport experiments, the quantum nature of matter becomes directly evident when changes in conductance occur only in discrete steps7, with a size determined solely by Planck’s constant $h$. Observations of quantized steps in electrical conductance7–9 have provided important insights into the physics of mesoscopic systems11 and have allowed the development of quantum electronic devices7. Even though quantized conductance should not rely on the presence of electric charges, it has never been observed for neutral, massive particles7. In its most fundamental form, it requires a quantum-degenerate Fermi gas, a ballistic and adiabatic transport channel, and a constriction with dimensions comparable to the Fermi wavelength. Here we report the observation of quantized conductance in the transport of neutral atoms driven by a chemical potential bias. The atoms are in an ultra-ballistic regime, where their mean free path exceeds not only the size of the transport channel, but also the size of the entire system, including the atom reservoirs. We use high-resolution lithography to shape the atom reservoirs. We use high-resolution lithography to shape light potentials that realize either a quantum point contact or a quantum wire for atoms. These constrictions are imprinted on a quasi-two-dimensional ballistic channel connecting the reservoirs7. By varying either a gate potential or the transverse confinement of the constrictions, we observe distinct plateaux in the atom conductance. The conductance in the first plateau is found to be equal to the universal conductance quantum, $1/h$. We use Landauer’s formula to model our results and find good agreement for low gate potentials, with all parameters determined a priori. Our experiment lets us investigate quantum conductors with wide control not only over the channel geometry, but also over the reservoir properties, such as interaction strength, size and thermalization rate.

As pointed out by Landauer in 1957, conduction is the transmission of carriers from one terminal to another8,9. At the heart of this idea is the separation of a conducting channel region from the terminals, which have the role of carrier reservoirs. Inelastic processes in these reservoirs keep them in a steady state, and a fast and dissipative dynamics ensures incoherent emission of the carriers into the channel. A criterion for the observation of quantized conductance is the suppression of both inelastic and elastic scattering processes in the channel, as well as an adiabatic connection of the channel to the reservoirs (Methods). These criteria are met in only very few systems, such as high-mobility two-dimensional (2D) electron gases10,11, atomic-sized break junctions in metals12,13, and carbon nanotubes11. The conducting channel is then described by a discrete set of modes in the transverse directions and a continuum along the transport direction. Each transverse mode contributes a maximum value of $1/h$ to the conductance $G$, reflecting the maximum phase-space occupation allowed by the Pauli principle. This gives

$$G = \frac{1}{h} \sum f(E_n - \mu)$$

where $f$ is the Fermi–Dirac distribution, $E_n$ is the energy of the $n$th transverse mode and $\mu$ is the chemical potential in the unbiased reservoirs. When the temperature is sufficiently low compared with the energy spacing of the transverse modes, the contribution of individual modes can be isolated in a transport measurement, leading to quantized plateaux in the conductance.

The composite nature of the Landauer picture of transport has long prevented its realization with cold atoms14, even though one-dimensional systems have been extensively investigated15–18. A situation opposite to the Landauer picture had been realized by placing Bose–Einstein condensates into a double–well potential, where the transport is governed by the coherence between the reservoirs19. Fermion reservoirs connected to 2D channels have recently been realized, but so far it has been possible to study transport only in the semi-classical regime, where individual transport modes are not resolved20,21. In sharp contrast to their solid-state counterparts, the atomic reservoirs are isolated systems, in which energy and particle number are strictly conserved and the dynamics is governed by free-particle motion interrupted only by rare elastic collisions between the particles.

The basis of our experiment is the transport set-up7. In brief, a weakly interacting gas of $N = 7.5(3) \times 10^4$ fermionic $^6$Li atoms is prepared in a cigar-shaped trap (the number in parentheses is the uncertainty (s.d.) in the final digit), which is then split into two reservoirs connected by a 2D channel using the repulsive potential of a TEM01-like (transverse electromagnetic) mode of a laser operating at 532 nm (Fig. 1a). The Fermi temperature in this trapping geometry is $T_F = 385(12)$ nK and the temperature is $T = 42(8)$ nK. A quantum point contact22 (QPC) is created by lithographically projecting a split-gate structure onto the 2D channel (Fig. 1a, b). To do so, the negative of a slit with a width of 12 μm is printed on a binary mask and illuminated with a laser beam at 532 nm. A projection system, consisting of an achromatic lens and a high-numerical-aperture microscope objective, shrinks the object by a factor of 11. The width of the resulting QPC in the channel region is measured to be 1.5(3) μm (full-width at half-maximum) in the $x$ direction, using a second, identical microscope placed opposite the first23. Because this width is comparable to the Fermi wavelength of 2.0 μm in our system, a single-mode regime should be accessible. The finite width of the diffraction-limited point spread function of the projection system leads to a harmonically confining potential along the $x$ axis.

The QPC is characterized by $v_x$ and $v_y$, the trap frequencies along $x$ and $y$ at the centre of the QPC, which originate from the harmonic confinement of the TEM01-like laser beam and the lithographically imprinted constriction, respectively. Typical values are $v_x = 10.0(4)$ kHz and $v_y = 30(3)$ kHz. The three lowest modes are thus separated by $\hbar v_x \approx 0.5$ μK, which is much more than the energy equivalent to the temperature of the gas. The zero-point energy in the QPC is $\hbar E_0 = (\hbar v_x + \hbar v_y)/2 \approx 1.0$ μK, which is larger than the chemical potential $\mu = 370(11)$ nK imposed by the reservoirs. We use an additional laser beam creating an attractive gate potential $V_g$ at the position of the QPC to successively populate its otherwise empty transverse modes (Fig. 1c–e). This laser beam has a wavelength of 767 nm and a waist of 25.0(6) μm, and is propagating along the $z$ axis.

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To measure the conductance, we prepare an initial particle number imbalance $\Delta N_0 = (N_L - N_R)_0 = 0.40(2)N$ between the two reservoirs, with $N_L$ and $N_R$ denoting the particle numbers in the left and right reservoirs, respectively. This leads to a chemical potential bias $\Delta \mu = 94(7) \text{nK} \ll h v_2$ driving a current $I = G \Delta \mu$ across the QPC. We access the conductance by measuring the relative particle number imbalance after 1.5 s of transport, assuming a linear response, and evaluating the compressibility of the reservoirs (Methods).

Figure 2 presents the measured conductance as a function of $V_g$ for two different vertical confinement frequencies, $v_z = 10.4 \text{kHz}$ and $v_z = 8.2 \text{kHz}$. The confinement frequency along $x$ is set to $v_x = 31.8 \text{kHz}$. Both curves start at zero conductance (note that the 8.2 kHz curve is artificially offset for clarity) because the QPC is entirely closed at small gate potentials owing to its zero-point energy. The conductance starts to rise at the point where $V_g$ compensates the zero-point energy, and saturates very close to the universal value $1/h$ as soon as the ground-state mode is tuned below the chemical potential of the reservoirs. Higher modes follow accordingly on further increasing $V_g$. We clearly resolve the first two conductance plateaux in the case of $v_z = 10.4 \text{kHz}$ (open blue circles), corresponding to the population of the $(n_x = 0; n_z = 0, 1)$ modes, where $n_x$ and $n_z$ are the quantum numbers in the harmonic potentials in the $x$ and $z$ directions. For $v_z = 8.2 \text{kHz}$ (filled red squares), the plateaux are narrower because the modes are more closely spaced in energy. In this case, we resolve the first three modes $(n_x = 0; n_z = 0, 1, 2)$ and even the onset of the $(0, 3)$ mode is visible.

The inset in Fig. 2 shows a close-up view of the first conductance plateau (with no offset). When $V_g$ is replaced with the total energy $E_{tot} = V_g + \mu$ of the particles minus $E_0$ and this value is normalized to $h v_2$, the two data sets fall on top of each other as a consequence of universality, with the width of the plateau given by $h v_2$. The absolute accuracy of our conductance measurement is limited by the uncertainty in the compressibility of the reservoirs, which amounts to 11% (Methods).

The quantization of conductance is universal and should not depend on the control parameter. To demonstrate this, we next use the horizontal confinement frequency $v_y$ of the QPC as a tuning parameter and keep the gate potential fixed. This is the counterpart of the measurement in solid-state physics where the split-gate voltage is tuned to reveal quantized conductance$^{2,3}$. The blue data points in Fig. 3 present this measurement with the horizontal cut is an image of the projected split-gate structure. The green solid line and red dashed line are $1/e^2$ contours of the intensity profiles of the laser beams creating the split-gate structure and the gate potential, respectively. d, Effective potential along the transport axis, consisting of the transverse-mode energy and the gate potential (Methods). The energy levels corresponding to the three lowest transverse modes are labelled $E_0$, $E_1$, and $E_2$. The depicted situation corresponds to the first plateau of the $v_x = 10.4 \text{kHz}$ data of Fig. 2. e, Energy dispersion relation of the particles in the QPC. The parabolas are offset by the quantized transverse-mode energies. Left and right reservoirs are represented by grey boxes, having chemical potentials $\mu \pm \Delta \mu/2$. 

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**Figure 1** | An atomic QPC. a, Lithographic imprinting of the QPC (vertical green beam). An achromatic lens and a high-numerical-aperture microscope objective are used to shrink the QPC structure onto the 2D channel region in the atomic cloud. An attractive gate potential is created by a red detuned laser beam (red beam). It is combined with the green beam on a dichroic mirror and focused onto the centre of the QPC. The TEM$_{00}$-like laser mode, which creates the 2D channel and sets the confinement frequency $v_y$ of the QPC, is shown as a horizontal green beam propagating along the $–x$ axis. b, Close-up view of the channel region: reservoirs, 2D channel and QPC are smoothly connected to each other. c, Imaged atomic density in the QPC for $v_y = 4.6 \text{kHz}$, $v_x = 29 \text{kHz}$ and $V_g = 1.4 \mu\text{K}$ in grey scale. The overlaid elliptical green region

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**Figure 2** | Conductance as a function of gate potential. Open blue circles correspond to a vertical confinement frequency of $v_z = 10.4 \text{kHz}$. Filled red squares correspond to $v_z = 8.2 \text{kHz}$ and are vertically shifted by two units for clarity. Each data point represents the mean of six measurements, and error bars indicate one standard deviation. Solid lines are theoretical predictions based on the Landauer formula of conductance. The shaded regions reflect the uncertainties in the input parameters (see text). Dashed lines are continuations of the solid lines and correspond to a change in the effective potential (Methods). Inset, first conductance plateau as a function of reduced energy, showing universal scaling. Vertical dashed lines indicate the width of the first plateau, and the horizontal dashed line indicates the universal conductance value $1/h$. 

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for $v_z = 10.9 \text{kHz}$ and $V_g = 1.0(1) \text{\mu K}$. We identify two increasingly wider plateaux centred at $v_z = 20$ and $35 \text{kHz}$. They correspond to the presence of two or one conductance modes, respectively. The bending at $v_z = 12 \text{kHz}$ is due to the population of a third mode, the $(1,0)$ mode. When widening the QPC further, the conductance rises quickly because the next transverse energy levels $(n_x, n_z)$ are energetically very close. The red data points in Fig. 3 correspond to $v_z = 9.2 \text{kHz}$ and $V_g = 0.8(1) \text{\mu K}$, and show the same features as the blue data, with a reduced plateau width owing to a smaller $v_z$.

We compare our data with theoretical predictions (solid lines in Figs 2 and 3) of the Landauer formula in the limit of entirely adiabatic and ballistic transport (equation (1)) without any fit parameters (Methods). The input parameters $T$, $\mu$, $\Delta \mu$, $v_z$ and $V_g$ are all independently measured quantities. The positions and widths of the plateaux are in general well predicted. The conductances on the plateaux reach the universal values for moderate gate potentials. For larger gate potentials (Fig. 2), the measured conductance is lower than predicted. Possible reasons could be the presence of additional contact resistances in the 2D tapered regions connecting the reservoirs with the QPC (Methods and Extended Data Fig. 4), or a small non-adiabaticity in the motion of the particles introduced by the gate potential, which accelerates the particles more and more towards the QPC as its strength is increased.

The sharpness of the transition from one plateau to the next is set by the finite temperature. We checked numerically within the adiabatic approximation that broadening due to tunnelling below the barrier or reflections above it is much smaller than the relevant thermal broadening of the Fermi edge of $\sim 4k_B T$ (Ref. 5; $k_B$, Boltzmann’s constant). Furthermore, we do not observe any nonlinear effects due to the applied finite bias $\Delta \mu$ because it is also smaller than $\sim 4k_B T$.

Even with isolated reservoirs, which cannot dissipate energy, our QPC features quantized conductance that can be described by Landauer’s theory. This is remarkable because our reservoirs are close to a collisionless regime, with a collision mean free path of approximately 12 mm, which is about 40 times larger than the size of the cloud and corresponds to a scattering time of about 400 ms (Methods and Extended Data Fig. 1). We attribute this to the three-dimensional nature and anharmonicity of the trapping potential, which should prevent closed particles orbits when

**Figure 3 | Conductance as a function of horizontal confinement.** Open blue circles correspond to a vertical confinement frequency of $v_z = 10.9 \text{kHz}$ and a gate potential of $V_g = 1.0(1) \text{\mu K}$. Filled red squares correspond to $v_z = 9.2 \text{kHz}$ and $V_g = 0.8(1) \text{\mu K}$, and are vertically shifted by two units for clarity. Solid lines are theoretical predictions based on the Landauer formula of conductance. The shaded regions reflect the uncertainties in the input parameters (see text). Data averaging and error bars are the same as in Fig. 2.

**Figure 4 | Quantum wire: conductance as a function of gate potential.** Open black circles correspond to a vertical confinement frequency of $v_z = 8.7 \text{kHz}$, and filled green squares to a lower confinement frequency of $v_z = 5.5 \text{kHz}$. For the stronger confinement, clear conductance plateaux are observed. The weaker confinement illustrates the onset of the contribution from the first excited mode along the x direction (see text). Solid lines are theoretical predictions based on the Landauer formula of conductance. Shaded regions, data averaging and error bars are the same as in Fig. 2. The inset shows the lithographically imprinted wire as imaged by a second, identical microscope objective. Below it, the corresponding effective potentials are drawn for the three lowest transverse modes, $E_0$, $E_1$ and $E_2$, for the parameters at the first plateau of the $v_z = 8.7 \text{kHz}$ data.
are observed for tight confinement ($v_1 = 24.6$ kHz and $v_2 = 8.7$ kHz; black circles in Fig. 4), using the gate potential as a tuning parameter. They are well reproduced by theory, up to a 10–20% reduction in height for the second and third plateaux. Possible interference effects, such as above-barrier resonances27, are not observed, probably because the finite temperature leads to an atomic coherence length of $\gamma(1)$ μm, which is less than the length of the quantum wire.

The green data in Fig. 4 demonstrate how the plateaux disappear when $v_1$ is lowered to 5.5 kHz. The conductance starts to increase at a lower gate potential because of the reduced zero-point energy. Well-defined plateaux for tight confinement ($v_1 \approx 1.75$ μK, which is where the first excited mode along the x direction begins to contribute. More precisely, when $v_1$ is increased from 0.25 to 1.75 μK the $(n_x = 0; n_z = 0, 1, 2, 3)$ modes begin to contribute, whereas above $V_g = 1.75$ μK the $(n_x = 0, n_z = 4, 5, \ldots)$ modes and the $(n_x = 1, n_z = 0, 1, \ldots)$ modes begin to contribute, giving rise to a change in slope by a factor of ~2. This feature is accurately predicted by theory.

With our projection technique, channels shorter than the Fermi wave-length can be straightforwardly implemented to study the breakdown of the adiabatic assumption and to produce structured tunnel contacts as building blocks for quantum dots or interferometers. Increasing the interactions should allow us to investigate a possible breakdown of quantized conductance and the Landauer picture in the strongly correlated regime29. The detection of conductance at the level of single quanta provides access to the physics of topological edge states30, transport in the vicinity of a quantum phase transition31, and universal conductance fluctuations32.

Online Content Methods, along with any additional Extended Data display items and Source Data, are available in the online version of the paper; references unique to these sections appear only in the online paper.

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1. Imry, Y. in Directions in Condensed Matter (eds Grinstein, G. & Mazenko, G.) 120–145 (World Scientific, 1986).
2. van Wees, B. J. et al Quantized conductance of point contacts in a two-dimensional electron gas. Phys. Rev. Lett. 60, 848–850 (1988).
3. Wharam, D. A. et al One-dimensional transport and the quantisation of the ballistic resistance. J. Phys. C 21, L209 (1988).
4. Imry, Y. Introduction to Mesoscopic Physics (Oxford Univ. Press, 2002).
5. Ihn, T. Semiconductor Nanostructures (Oxford Univ. Press, 2010).
6. Sato, Y., Eom, B.-H. & Packard, R. On the feasibility of detecting quantized conductance in neutral matter. J. Low Temp. Phys. 141, 99–109 (2005).
7. Brandt, J.-P., Meineke, J., Stadler, D., Krinner, S. & Esslinger, T. Conduction of ultracold fermions through a mesoscopic channel. Science 337, 1069–1071 (2012).
8. Landauer, R. Spatial variation of currents and fields due to localized scatterers in metallic conduction. IBM J. Res. Develop. 1, 223–231 (1957).
9. Büttiker, M., Imry, Y., Landauer, R. & Pinhas, S. Generalized many-channel conductance formula with application to small rings. Phys. Rev. B 31, 6207–6215 (1985).
10. Krans, J. M., van Ruitenbeek, J. M., Fisun, V. V., Yanson, I. K. & de Jongh, L. J. The signature of conductance quantization in metallic point contacts. Nature 375, 767–769 (1995).
11. Frank, S., Poncharal, P., Wang, Z. L. & de Heer, W. A. Carbon nanotube quantum resistors. Science 280, 1744–1746 (1998).
12. Thywissen, J. H., Westervelt, R. M. & Prettis, M. Quantum point contacts for neutral atoms. Phys. Rev. Lett. 83, 3762–3765 (1999).
13. Gorlitz, A. et al. Realization of Bose-Einstein condensates in lower dimensions. Phys. Rev. Lett. 87, 130402 (2001).
14. Moritz, H., Stoferle, T., Kühn, M. & Esslinger, T. Exciting collective oscillations in a trapped 10 gas. Phys. Rev. Lett. 91, 250402 (2003).
15. Paredes, B. et al. Tonks–Girardeau gas of ultracold atoms in an optical lattice. Nature 429, 277–281 (2004).
16. Kinoshita, T., Wenger, T. & Weiss, D. S. Observation of a one-dimensional Tonks–Girardeau gas. Science 305, 1125–1128 (2004).
17. Bouchoule, L., van Druten, N. & Westbrook, C. in Atom Chips (eds Reichel, J. & Vuletic, V.) 331–363 (Wiley, 2010).
18. Serwane, F. et al. Deterministic preparation of a tunable few-fermion system. Science 332, 336–338 (2011).
19. Albiez, M. et al. Direct observation of tunneling and nonlinear self-trapping in a single bosonic Josephson junction. Phys. Rev. Lett. 95, 010402 (2005).
20. Stadler, D., Krinner, S., Meineke, J., Brandt, J.-P. & Esslinger, T. Observing the drop of resistance in the flow of a superfluid Fermi gas. Nature 491, 736–739 (2012).
21. Krinner, S., Stadler, D., Meineke, J., Brandt, J.-P. & Esslinger, T. Superfluidity with disorder in a thin film of quantum gas. Phys. Rev. Lett. 110, 100601 (2013).
22. van Houten, H. & Beenakker, C. Quantum point contacts. Phys. Today 49, 22–27 (1996).
23. Zimmermann, B., Müller, T., Meineke, J., Esslinger, T. & Moritz, H. High-resolution imaging of ultracold fermions in microscopically tailored optical potentials. New J. Phys. 13, 043007 (2011).
24. Glazman, L. I. & Lesovik, G. B. Khmel’nit’skii, D. E. & Shekhter, R. I. Reflectionless quantum transport and fundamental ballistic-resistance steps in microscopic constrictions. JETP Lett. 48, 238–241 (1988).
25. Kouwenhoven, L. P. et al. Nonlinear conductance of quantum point contacts. Phys. Rev. B 39, 8040–8043 (1989).
26. Ulriceh, S. & Zwerg, W. Where is the potential drop in a quantum point contact? Superlattices Microstruct. 23, 719–730 (1998).
27. Zsafer, A. & Stone, A. D. Theory of quantum conduction through a constriction. Phys. Rev. Lett. 62, 300–303 (1989).
28. Yacoby, A. & Imry, Y. Quantization of the conductance of ballistic point contacts beyond the adiabatic approximation. Phys. Rev. B 41, 5341–5350 (1990).
29. Vignale, G. & Di Ventra, M. Incompleteness of the Landauer formula for electronic transport. Phys. Rev. Lett. 80, 1748–1751 (1998).
30. Hasan, M. Z. & Kane, C. L. Topological insulators. Rev. Mod. Phys. 82, 3045–3067 (2010).
31. Sachdev, S. Quantum Phase Transitions 260–290 (Cambridge Univ. Press, 2011).
32. Lee, P. A. & Stone, A. D. Universal conductance fluctuations in metals. Phys. Rev. Lett. 55, 1622 (1985).
METHODS

Experimental set-up. Our experiment uses the system described in refs 7, 33. In brief, a degenerate Fermi gas of $^6$Li atoms is produced by evaporative cooling of a balanced spin-mixture of the lowest and third-lowest hyperfine states in a hybrid magnetic and optical dipole trap in a homogeneous magnetic field of 388 G, using a magnetic field gradient. The dipole trap operates at a wavelength of 1.064 nm, and has a waist of 70 μm and a trap depth of 1.2 μK. The experiments are performed in a homogeneous magnetic field of 552 G, where the scattering length is $-187 a_0$, $a_0$ denoting the Bohr radius. This allows for sufficiently fast thermalization of the reservoirs, and at the same time ensures a ballistic transport channel. All quantities in the text are stated for a single hyperfine state, implying that the conductance plateaux appear in multiples of $1/h$ and not $2/h$.

Transport sequence. During the preparation of the degenerate gas, a magnetic field gradient of 0.2 mT m$^{-1}$ is applied along the $y$ axis to shift the trap with respect to the QPC. A repulsive elliptic gate beam focused on the centre of the QPC is used to separate the two reservoirs as in ref. 33. The powers of the laser beams creating the 2D channel and the QPC are consecutively ramped from zero to their final values within 200 ms. Then evaporative cooling is enforced by using a magnetic field gradient along the $z$ axis. This procedure results in a well-defined particle number imbalance between the two reservoirs. Next the dipole trap is adiabatically decompressed from a trap depth of 5.6 μK to a final depth of 1.2 μK within 200 ms to further reduce the absolute temperature of the gas. During the same time interval the attractive gate potential is ramped from zero to $V_g$. Finally, the magnetic field gradient along $y$ is ramped to zero within 60 ms, resulting in a well-defined chemical potential difference $\Delta \mu$ between the two reservoirs. We start the transport process by removing the repulsive gate beam. After a transport time of 1.5 s we switch the repulsive gate beam back on to stop the transport process, and we measure the atom number in both reservoirs via absorption imaging along the $x$ direction.

Conductance evaluation. From the equation for the current, $I = G \Delta \mu$, we find that the temporal evolution of the particle number imbalance between the left and right reservoirs, $\Delta N(t) = N_l(t) - N_r(t)$, is governed in linear response by the equation

$$\frac{d}{dt} \Delta N(t) = - \frac{G}{C_{\text{eff}}}(\Delta N(t))$$

where $N_l(t) + N_r(t)$ is the total atom number and $C_{\text{eff}} = (1/\rho_0 + 1/\rho_1)^{-1}$ is an effective compressibility determined by the respective compressibilities of the single reservoirs at equilibrium. In agreement with this equation we observe an exponential decay of the relative particle number imbalance as a function of time, the time constant being $\tau = C_{\text{eff}}/G$. We evaluate $G$ by measuring $\tau$ and evaluating $C_{\text{eff}}$. The error made when replacing $\Delta \mu$ with $\Delta N(t)/\Delta N(0)$ to obtain the linear response equation (equation (2)) is smaller than $5\%$ at the largest value of $\Delta N(0)$, which is 0.4.

We determine $\tau$ by measuring the relative particle number imbalance at $t = 0$ and again after a transport time of $\tau = 1.5$ s. From the solution of equation (2) we obtain

$$\frac{1}{\tau} = \frac{1}{\tau_0} \log \left(\frac{\Delta N(t)}{\Delta N(t_0)}\right)$$

where a constant offset of $-0.2(1)$ is subtracted from both balances. This offset originates partly from the relative alignment of the gate beam and the lithographic system, and for the data of Fig. 2 it varied linearly from $-0.2$ to $-0.16$ over the range shown.

The compressibilities $C_0 = C_\infty = C$ of the identical reservoirs are calculated from the trap geometry, particle number and temperature, assuming a non-interacting Fermi gas. In brief, the trapping potential is harmonic along the $x$ and $z$ directions and half-harmonic along the $y$ direction, with respective trapping frequencies of 194, 235 and 157 Hz along the $x$, $y$ and $z$ directions. The effect of the repulsive potential of the TEM$_{10}$-like laser mode, creating the 2D channel, is to shift the Fermi energy and chemical potential by 17% towards larger values with respect to the unperturbed cloud. The compressibility is almost not affected. The systematic uncertainty in $C_{\text{eff}}$ is 11%, which is due to the calibration error in the total particle number and an uncertainty in the overall trapping potential.

Trapping frequencies of the QPC and the quantum wire. The transverse trapping frequencies of the QPC and quantum wire, $v_x$ and $v_y$, are measured by parametric heating in a dipole trap created by a laser beam with a waist of 8 μm and a wavelength of 767 nm propagating along the $z$ axis. The observed resonances have relative widths (full-widths at half-maximum) of $\Delta v_x/v_x = 0.05$ and $\Delta v_y/v_y = 0.30(5)$, respectively. $v_y$ is found to depend weakly on $v_x$. This is because the darkness of the projected QPC structure decreases as a result of diffraction when moving out of focus along the $z$ axis, thus creating an additional confinement along $z$. We measure this contribution to be $v_z = 0.16 v_x$. Hence, $v_y$ is given by $v_y = \sqrt{v_x^2 + 0.16 v_x^2}$, where $v_y$ is the trapping frequency in the absence of the QPC. The values of $v_y$, stated in the text for the data of Fig. 3, are evaluated for $v_x = 31.8$ kHz to be comparable to the values set for the data of Fig. 2. For the shown range of $v_y \in [10 \text{ Hz}, 50 \text{ Hz}]$, $v_y$ varies by 12% around its mean value for the $v_x = 10.9 \text{ kHz}$ data and by 16% around its mean value for the $v_x = 9.2 \text{ kHz}$ data.

Adiabatic approximation and theory curves. The computation of the conductance makes use of the adiabatic approximation, allowing for a separation of longitudinal ($y$) and transverse ($x,z$) variables. It neglects scattering between different transverse modes and is justified if the confinement of the constriction varies smoothly along the transport direction. This is to a good approximation the case for both the QPC, with its Gaussian envelope, and the quantum wire, with its triangular openings that are smoothened by the inherent low-pass filtering of the projection system. In the resulting one-dimensional Schrödinger equation the transverse energy $E_n(y) = E_{n_z}(y) + hv_x f_x(y) (n_x + 1/2) + hv_y f_y(y) (n_y + 1/2)$, with $f_x(y)$ and $f_y(y)$ describing the spatial variations of the corresponding trapping frequencies, acts as an additional potential. Together with the gate potential $V_g(y)$ it forms the effective potential drawn in Fig. 1d and in the inset of Fig. 4. For the QPC we have $f_x(y) = \exp(-y^2/\omega_x^2)$, with $\omega_x = 5.6(3) \mu$m and $\omega_y = 30(1) \mu$m, whereas for the quantum wire $f_y$ is constant along the wire, with its edges modelled by an error function.

The theory lines in Fig. 2 are dashed above $V_g = 1.5 \mu$m because at this point the maximum of the effective potential moves from the centre to the sides of the QPC, which is not taken into account in the theory. This effect is not expected to explain the observed shift of the conductance below the universal values for large gate potentials.

The conductance is calculated from equation (1), which we obtain from the two-terminal Landauer formula in the adiabatic regime

$$G = \frac{1}{h} \sum n f \left( E_n - \frac{\partial f}{\partial E} (E - \mu) \right)$$

By setting the transmission probability to be $T(E) = \Theta(E - E_n)$, with $\Theta(E - E_n)$ the Heaviside step function. This substitution corresponds to the semiclassical approximation, where the transmission probability for particles is 1 if their total energy is larger than their transverse energy, and 0 otherwise. It neglects tunnelling below the barrier and reflections above it$^{22}$, which would lead to a broadening of $T(E)$.

Taking into account the gate potential for the conductance evaluation, equation (1) reads

$$G = \frac{1}{h} \sum n f \left( E_n - V_g \frac{\partial f}{\partial E} (E - \mu) \right)$$

In the case of the quantum wire $V_g$ is multiplied by a factor of $\exp(-2 \times 9.5^2/25^2)$ to account for the fact that the points of largest effective potential are located $9.5 \mu$m from the centre of the gate potential and the quantum wire. The shaded error regions in Figs 2–4 cover statistical and systematic errors of the input parameters and are determined using Gaussian error propagation. The main contribution in Figs 2 and 4 is the uncertainty in $v_n$, whereas in Fig. 3 it is the uncertainty in $v_g$. Timescales in the reservoirs. Here we discuss the timescales and associated length scales relevant to the evolution of the coupled reservoirs, as sketched in Extended Data Fig. 1.

The shortest timescale is the Fermi time $h/k_B T_F = 125 \mu$s. On this timescale the atoms exiting the channel into the reservoirs travel a distance of the order of the Fermi wavelength. Recent calculations in a geometry similar to our experimental configuration, where a one-dimensional channel enters a wider region, show that it takes a time of only $10h/k_BT_F$ for a steady flow pattern to establish$^{16}$. Importantly, it was shown that this was the case in the presence of an incoherent bath as well as in a complete microcanonical picture.

The next-longest timescale is the coherence time $h/k_BT = 1.1$ ms. This timescale is about an order of magnitude smaller than the oscillation period in the reservoirs $1/\nu = 8.5$ ms, where $\nu$ is the mean trap frequency in the reservoirs. This prevents interference patterns from developing in the reservoirs. The dephasing of the interferes is due to thermal averaging, and not due to inelastic scattering with phonons, as can happen in solid-state systems.

Because the reservoirs are three-dimensional and the trapping potential is anharmonic, it is unlikely that trajectories will be closed, and we should instead expect complex orbits leading to escape times much larger than the oscillation period$^{24}$. The thermalization of the incident particles from the QPC occurs on the scale of the scattering time $\tau_s = (1/\nu)(T/T_F) = 400$ ms, where $\nu$ is the peak density at the trap centre, $\sigma$ is the scattering cross-section for interparticle collisions and $\nu$ is the Fermi velocity. This timescale does not account for the finite bias, and should thus be considered an upper bound. Here we are concerned with the decay of excitations above the Fermi surface and not with a generic particle in the trap, and the Pauli principle thus leads to suppression of scattering by a factor $T/T_F$ and not $(T/T_F)^2$. This situation is similar to the case of evaporative cooling where the exit states are empty$^{27}$. Note that this scattering is elastic, and that the many-body evolution of
the reservoir remains coherent. The interparticle scattering ensures that the energy distribution follows Fermi–Dirac statistics. While theoretically possible, revivals due to the coherent many-body evolution of the reservoirs should take place only on extremely long timescales. The corresponding collision mean free path is about 12 nm, showing that many oscillations in the reservoirs take place before collisions redistribute the energy. This also shows that in our experiment, the decay of excitations introduced by the transport is non-local.

The timescale for current flow corresponds to the timescale $t_{\text{RC}}$ of the equivalent RC circuit. In our case, $t_{\text{RC}} = \tau_{\text{qpc}} = 8\ s$ for the first conductance plateau. We measure the decay of the atom number imbalance in the reservoirs after an observation time of typically $t_{\text{obs}} = 1.5\ s$.

Inelastic processes take place even on longer timescales. Three-body recombinations are strongly inhibited by the Pauli principle. All the available estimates for the loss coefficient in the BEC–BCS crossover (with much stronger interactions than in our present work) suggest a timescale larger than 1,000 s (ref. 38). Spontaneous emission in the far-detuned optical dipole trap takes place on a timescale of $t_{\text{trap}} = 580\ s$. The attractive gate potential has a spontaneous emission timescale of $t_{\text{gate}} = 63\ s$. Yet only a tiny fraction of the atoms are exposed to this potential (<1%), which reduces accordingly the actual probability for an atom to scatter from this beam. The finite depth of the optical dipole trap allows in principle for free evaporation of high-energy atoms. This effect is controlled by an effective truncation parameter $\kappa = (U - k_{q}T_{q})/k_{B}T = 19$, which is large enough that escape of atoms can be neglected.

Non-interacting reservoirs. We took another set of data with parameters similar to the ones of Fig. 2, with the interaction strength tuned to zero. The reservoirs are prepared as described in the main text, with the same initial particle numbers and a temperature of $T = 50(10)\ \text{mK} = 0.13T$. The confinement in the QPC was $v_{z} = 10.4\ \text{kHz}$ and $v_{x} = 31.8\ \text{kHz}$, and the observation time was set to 4 s. Before starting the transport process, the magnetic field was ramped to 587 G, where the scattering length amounts to $a = -4a_{0}$ (ref. 39). This corresponds to an elastic scattering time of 850 s for atoms above the Fermi level, placing the reservoirs in non-interacting conditions. Even though the cloud is prepared in thermal equilibrium, the subsequent evolution of the reservoirs is non-thermal. Applying the same data processing as for the weakly interacting gas, we obtain the same step features, as presented in Extended Data Fig. 2 (filled magenta squares). The data was taken simultaneously with a reference data set (open blue circles in Extended Data Fig. 2), where the scattering length was set to $a = -187a_{0}$, as it was for all the data presented in the main text. By comparison, we see that the complete absence of interactions does not lead to any noticeable change in the condensate. The theory curve in Extended Data Fig. 2 includes a common envelope of the QPC confinement, $f_{y}(x) = e^{-x^{2}/w^{2}}$ (see above), equation (3) reads

$$\left| \frac{\kappa_{y}}{\Delta_{y}} \right| \ll \Delta_{y}$$

Because $\Delta_{y}$ is very small far away from the QPC, this criterion can be fulfilled only in the region of the QPC and necessarily fails at its entrance and exit. The criterion to observe quantized conductance relies on the fact that no transitions between transverse modes take place inside the point contact. In principle, this could be the case because in their rest frame, the moving atoms see a time-dependent confinement leading to possible transitions between modes when this change is too fast. Specifically, atoms will stay in the same transverse energy state $\mathcal{E}_{y}$, relative to the level spacing $\Delta = \hbar \Delta_{y}$ is small compared to $\Delta_{y}$ itself:

$$\left| \frac{\kappa_{y}}{\Delta_{y}} \right| \ll \Delta_{y}$$

$$\mathcal{E}_{y} = \hbar \omega_{0}, \quad \text{if in the frame moving with the atoms, the temporal change of this energy } \hbar \omega_{0}, \text{ relative to the level spacing } \Delta = \hbar \Delta_{y} \text{ is small compared to } \Delta_{y} \text{ itself} \text{[40]}.$$
and the regime where the resistances of subsequent constrictions add, in which case $G$ would be given by a formula interpolating between the two regimes$^{41}$. Furthermore, the latter regime might be reached only for large gate potentials, in which case equation (5) should be applied only for large gate potentials. However, for small gate potentials, that is, for small conductances $G$, the contribution of $G_{2D}$ in the evaluation of $G_{QPC}$ is anyway very small, because the resistor $1/G_{QPC}$ then dominates the two resistors $1/G_{2D}$.

33. Brantut, J.-P. et al. A thermoelectric heat engine with ultracold atoms. Science 342, 713–715 (2013).
34. Hung, C.-L., Zhang, X., Gemelke, N. & Chin, C. Accelerating evaporative cooling of atoms into Bose-Einstein condensation in optical traps. Phys. Rev. A 78, 011604 (2008).
35. Beria, M., Iqbal, Y., Di Ventra, M. & Müller, M. Quantum-statistics-induced flow patterns in driven ideal Fermi gases. Phys. Rev. A 88, 043611 (2013).
36. Gattobigio, G. L., Couvert, A., Georgeot, B. & Guéry-Odelin, D. Exploring classically chaotic potentials with a matter wave quantum probe. Phys. Rev. Lett. 107, 254104 (2011).
37. O’Hara, K. M., Gehm, M. E., Granade, S. R. & Thomas, J. E. Scaling laws for evaporative cooling in time-dependent optical traps. Phys. Rev. A 64, 051403 (2001).
38. Du, X., Zhang, Y. & Thomas, J. E. Inelastic collisions of a fermi gas in the BEC-BCS crossover. Phys. Rev. Lett. 102, 250402 (2009).
39. Zürn, G. et al. Precise characterization of $^{6}$Li Feshbach resonances using trap-sideband-resolved RF spectroscopy of weakly bound molecules. Phys. Rev. Lett. 110, 135301 (2013).
40. Migdal, A. Qualitative Methods in Quantum Theory 115–118 (Perseus, 2000).
41. Büttiker, M. Role of quantum coherence in series resistors. Phys. Rev. B 33, 3020–3026 (1986).
42. Ulreich, S. Transport Durch Ballistische Leiter 44–52, PhD thesis, LMU München (1997).
Extended Data Figure 1 | Distribution of timescales in the reservoirs. The corresponding physical phenomena are indicated on the time line. The meanings of the different notations are defined in the text.
Extended Data Figure 2 | Conductance as a function of gate potential for non-interacting reservoirs. Filled magenta squares correspond to a scattering length of $a = -4a_0$, whereas open blue circles represent a reference data set for weakly interacting reservoirs, where $a = -187a_0$. Each data point is the mean of nine measurements, and error bars indicate one standard deviation. The black line is a theoretical prediction based on the Landauer formula of conductance, and the shaded region reflects the uncertainties in the input parameters.
Extended Data Figure 3 | Illustration of the adiabaticity criterion.
Transverse energy level spacing $\Delta \omega_x = \Delta E_x/h$ along the tightly confined QPC direction (solid red line), and corresponding temporal change in the ground-state trapping frequency in the moving frame of the atoms $\omega_{x,n_z=0}/\Delta \omega$ as a function of position along the QPC for the three possible values of $n_z$ when $q = 3$ modes are populated.
Extended Data Figure 4 | Conductance of the 2D region and of the QPC as a function of gate potential. 

a, Conductance $G_{2D}$ as a function of gate potential in the absence of the QPC. The solid line is a linear fit to the data.

b, Conductance $G_{QPC}$ of the QPC only, when considering the contact resistances of the 2D confinement as series resistors to the QPC. The data set, colour code, solid line and shaded region are the same as in Extended Data Fig. 2. Error bars are obtained by propagating the errors in $G$ and $G_{2D}$. 