the lattice sites
operators
dependence in the correlations, leading to:
For a regular three-dimensional system the structure term converges to:

\[ \frac{2}{\pi^{1/2} \sigma^{3/2}} \int_{BZ} C_{m}(k) \frac{\sigma}{x + 1} \mathrm{d}k = \left\{ \begin{array}{ll}
\frac{2}{\pi^{1/2} \sigma^{3/2}} \int_{BZ} C_{m}(k) \frac{\sigma}{x + 1} \mathrm{d}k & \text{for regular three-dimensional system,} \\
\frac{2}{\pi^{1/2} \sigma^{3/2}} \int_{BZ} C_{m}(k) \frac{\sigma}{x + 1} \mathrm{d}k & \text{for a regular one-dimensional lattice with on-site operator } n_l, \text{ the second term introduces a spatial dependence in the correlations, leading to:}
\end{array} \right. \]

\[ C_{m}(d) = C(x_1 - x_2) = 1 + \sum_{n, m = 0}^{\infty} \frac{\sigma}{x + 1} \frac{2}{\pi^{1/2} \sigma^{3/2}} \int_{BZ} C_{m}(k) \frac{\sigma}{x + 1} \mathrm{d}k \]

\[ C(d) = 1 + \frac{\sigma}{x + 1} \int_{BZ} C_{m}(k) \frac{\sigma}{x + 1} \mathrm{d}k \]

The heights of the peaks at the reciprocal lattice momenta therefore scale as \( N^{-1/2} \) for this simple homogeneous case. As indicated in the text, the \( N^{-1/2} \) scaling is modified to \( N^{-1/4} \) for our harmonically trapped system by the appearance of Mott domains with filling factor larger than one for higher atom numbers. The prediction for the exponent has been obtained by numerically evaluating the sum in equation (5) using a model distribution of atoms in the lattice sites confined by a global parabolic potential. This distribution is predicted assuming the system can be described in the strongly interacting limit with a local density approximation.

Received 27 January; accepted 28 February 2005; doi:10.1038/nature03500.


Acknowledgements We acknowledge discussions with E. Altman and M. Greiner, as well as financial support by the DFG, AFOSR and the EU under a Marie-Curie Fellowship (E.G.) and a Marie-Curie Excellence grant.

Competing interests statement The authors declare that they have no competing financial interests.

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Orbital Kondo effect in carbon nanotubes
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Progress in the fabrication of nanometre-scale electronic devices is opening new opportunities to uncover deeper aspects of the Kondo effect—a characteristic phenomenon in the physics of strongly correlated electrons. Artificial single-impurity Kondo systems have been realized in various nanostuctures, including semiconductor quantum dots, carbon nanotubes and individual molecules. The Kondo effect is usually regarded as a
spin-related phenomenon, namely the coherent exchange of the spin between a localized state and a Fermi sea of delocalized electrons. In principle, however, the role of the spin could be replaced by other degrees of freedom, such as an orbital quantum number. Here we show that the unique electronic structure of carbon nanotubes enables the observation of a purely orbital Kondo effect. We use a magnetic field to tune spin-polarized states into orbital degeneracy and conclude that the orbital quantum number is conserved during tunnelling. When orbital and spin degeneracies are present simultaneously, we observe a strongly enhanced Kondo effect, with a multiple splitting of the Kondo resonance at finite field and predicted to obey a so-called SU(4) symmetry.

The simplest Kondo system consists of a localized, spin-1/2 electron coupled to a Fermi sea by means of a Heisenberg-like exchange interaction. This simple system can be realized with a quantum dot (QD) device, which is a small electronic island connected to metallic leads via two tunnel barriers (Fig. 1a). Below a characteristic temperature $T_K$, the so-called Kondo temperature, a many-body singlet state is formed between the QD spin and the surrounding conduction electrons (Fig. 1a). This state adds a resonant level at the Fermi energy of the electrodes, enabling the tunnelling of electrons across the QD. Such a Kondo resonance can lead to a strong enhancement of the conductance, overcoming the Coulomb blockade effect. In principle, a Kondo effect might also occur in the absence of spin if another quantum number, for example an orbital degree of freedom, gives rise to a degeneracy (Fig. 1b). In this case, Kondo correlations lead to the screening of the local orbital ‘polarization’, and an orbital singlet is formed through a combination of orbital states. In the presence of both spin and orbital degeneracy, quantum fluctuations lead to a strong mixing of these two degrees of freedom (Fig. 1c). This increased degeneracy yields an enhancement of $T_K$ (ref. 11). In the low-temperature limit, this system is described by a Hamiltonian obeying SU(4) symmetry, which implies that the spin and charge degrees of freedom are fully entangled and the state of the electron is represented by a four-component ‘hypermultiplet’.

An orbital degeneracy is naturally expected in the electronic structure of carbon nanotubes (CNTs). This degeneracy can intuitively be viewed to originate from the two equivalent ways in which electrons can circle around the graphene cylinder, namely clockwise and anticlockwise. The rotational motion confers an orbital magnetic moment on the electrons. Consequently, the orbital degeneracy can be split by a magnetic field, $B$, parallel to the nanotube axis. (Experimental evidence for this effect, originally predicted by Ajiki and Ando, has recently been reported.) We label the orbital states of a CNT QD as $|+\rangle$ and $|--\rangle$ according to the sign of the energy shift that they experience under an applied $B$. Size quantization due to the finite CNT length results in two sets of orbital levels, $E_l^{(n)}$ and $E_l^{(2)}$, where $n = 1, 2, 3, \ldots$ is the quantization number. $E = E^{(n)}$ at $B = 0$ (assuming no orbital mixing), resulting in a fourfold degeneracy when including spin. The orbital and spin degeneracies are simultaneously lifted by a parallel $B$ (Fig. 1d). The use of $B$ permits the tuning of new degeneracies in connection with the crossing between levels from different shells. Here we are particularly interested in the crossing between states with the same spin polarization, of the type indicated by the yellow rectangle in Fig. 1d. We show below that the twofold orbital degeneracy originating from such a crossing gives rise to a purely orbital Kondo effect. We then consider the case of concomitant spin and orbital degeneracy at $B = 0$ (green rectangle in Fig. 1d) and present evidence for an SU(4) Kondo effect.

In a measurement of the linear conductance, $G$, as a function of gate voltage, $V_G$, the fourfold shell structure leads to consecutive groups of four closely spaced Coulomb blockade oscillations. The $B$-evolution of such oscillations is shown in Fig. 2a for a CNT QD device (described in the inset to Fig. 2a and in the figure legend) in a $V_G$ region encompassing two adjacent shells. Coulomb peaks (highlighted by green lines) appear as lines running from bottom to top and denote the sequential addition of electrons to the QD; the electron number increases from left to right. The observed pattern is explained in detail on the basis of the single-particle spectrum in Fig. 1d and taking into account the Coulomb interaction between electrons (see Supplementary Information).

The Coulomb peaks move to the left or right when increasing $B$,
which corresponds to adding the last electron to a $| - \rangle$ or $| + \rangle$ orbital, respectively. When the ground-state configuration of the QD changes, kinks appear in the $B$-evolution of the Coulomb peaks. The two enhanced-conductance ridges at $B = B_0 \sim 6 \, T$, bounded by two such kinks and highlighted by dotted yellow lines, are due to the crossing between $| - \rangle$ and $| + \rangle$ states as described in Fig. 1d. A detailed analysis (see Supplementary Information) indicates that along these ridges the QD ground state is doubly degenerate, with the last added electron occupying the level crossing between $| +, 1 \rangle$ and $| -, 1 \rangle$ (left ridge) or between $| +, 1 \rangle$ and $| -, 1 \rangle$ (right ridge).

In the region near the degeneracy point, we are able to measure a small coupling between orbital states\(^\text{21}\), resulting in level repulsion at $B = B_0$. The energy splitting is directly observed in the spectroscopy data of Fig. 2b, where the differential conductance, $dI/dV$, is shown as a function of $B$ and bias voltage, $V$. In this measurement $V_G$ and $B$ are simultaneously varied to follow the middle of the Coulomb valley (dashed blue line in Fig. 2a). Here, single-electron tunnelling is suppressed and the spectroscopy is based on higher-order co-tunnelling processes, which lead to an enhancement of $dI/dV$ every time that $V$ equals an internal excitation energy\(^\text{22}\). We focus on the high-transmission spin filter\(^\text{12–14}\). In contrast, the conductance enhancement that occurs for three-electron shell filling originates from Kondo correlations. The observed behaviour is characteristic of a split Kondo resonance—that is, a Kondo resonance associated with two quasi-degenerate states—in line with recent theoretical predictions\(^\text{24}\) and experiments\(^\text{25}\). It is important to note that the Zeeman spin splitting, $E_Z = g\mu_B B_0$, is three times $\delta_B$, indicating that the Kondo effect originates entirely from orbital correlations occurring at the crossing between two spin-polarized states, $| +, 1 \rangle$ and $| -, 1 \rangle$. This conclusion is in agreement with the zero-field data that we show below. The large Zeeman splitting also ensures that the observed orbital Kondo resonance provides a conducting channel only for $| \uparrow \rangle$ electrons, thereby acting as a high-transmission spin filter\(^\text{22–14}\). In contrast, the conductance enhancement that occurs for three-electron shell filling originates from $| +, 1 \rangle$ and $| -, 1 \rangle$ states and hence it allows only the tunnelling of $| \uparrow \rangle$ electrons. Switching from one degeneracy to the other is controlled by simply switching the gate voltage, which then causes the CNT QD to operate as a bipolar spin filter.

We now centre our attention on the zero-field regime, in which both orbital and spin degeneracies are expected (green rectangle in Fig. 1d). The Coulomb oscillations corresponding to the filling of a single shell are shown in Fig. 3a for a different CNT QD device. The four oscillations are clearly visible at 8 K (red trace). At lower $T$, the conductance exhibits a pronounced enhancement in regions I and III—that is, for 1 and 3 electrons on the shell—and the corresponding Coulomb blockade valleys disappear completely at 0.3 K (black trace). This conductance enhancement is a hallmark of Kondo correlations. From the $T$-dependence (fully shown in the Supplementary Information) we estimate that $T_K \approx 7.7 \, K$, an unusually high value that can be ascribed to the enhanced degeneracy\(^\text{21}\).

Figure 2 Orbital Kondo effect. a. Colour-scale representation of the linear conductance, $G$, versus $B$ and $V_G$ at $T = 30 \, mK$. $G$ increases from dark blue to dark red. The green lines highlight the $B$-evolution of the Coulomb peaks. The dotted yellow lines highlight regions of enhanced conductance due to Kondo effect. Roman labels indicate the number of electrons on the last occupied shell near $B = 0$. Orange numbers indicate the spin of the ground state. Inset: device scheme. Carbon nanotubes were grown by chemical vapour deposition on p-type silicon substrates with a surface oxide 250 nm thick. Individual nanotubes were located by atomic force microscopy and contacted with Ti/Au electrodes (typical separation $\sim 100–800 \, nm$) defined by e-beam lithography. The highly doped silicon substrate was used as a back-gate. b. Colour-scale plot of the differential conductance, $dI/dV$, against $V$ and $B$ along the dashed blue line in a. The field splits the Kondo resonance into multiple peaks. The two orange lines highlight the evolution of the peaks associated with the spin and orbital splitting, respectively. The spectroscopy features are more pronounced for $V < 0$, most probably due to asymmetric tunnel barriers\(^\text{20}\). The yellow lines highlight the orbital anticrossing at $B = B_3 = 5.9 \, T$. c. Coulomb diamond for one electron on the last occupied shell at $B = 5.9 \, T$. d. Plot of $dI/dV$ against $V$ with varying $T$, from 25 mK (thick blue trace) to 1.1 K (thick red trace), at the anticrossing point ($B = 5.9 \, T$, $V_G = 937 \, mV$). Orbital splitting, $\delta_B$, and Zeeman splitting, $E_Z$, are compared visually. The split Kondo peaks decrease with increasing $T$. Inset: plot of peak height against $T$ evaluated for the left peak.

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The important contribution of the orbital degree of freedom becomes apparent from the $B$-dependence of $G$ (Fig. 3c, f). If this Kondo effect were determined by spin only (this could be the case if one of the orbitals was coupled weakly to the leads), $G$ should decrease on a field scale $B \sim k_B T_K / g \mu_B \sim 6 \text{T}$ due to Zeeman splitting\cite{1}. In contrast, $G$ decays on a much smaller scale, $B \sim k_B T_K / 2 \mu_{\text{orb}} \sim 0.5 \text{T}$, which is determined by the orbital splitting (an estimate of the orbital magnetic moment, $\mu_{\text{orb}}$ is given below).

In the nonlinear regime, a single zero-bias Kondo resonance appears in regions I and III (Fig. 3b). Contrary to the result in Fig. 2c, no orbital splitting is observed because of the much larger $T_K$ ($k_B T_K > \delta$ (refs 12, 13, 21)). In region II we observe two peaks at finite bias, reflecting the already known splitting of a singlet–triplet Kondo resonance\cite{1}. To show that the Kondo resonance in regions I and III arises from simultaneous orbital and spin Kondo correlations, we investigate the effect of lifting spin and orbital degeneracies at finite $B$. As opposed to an ordinary spin-$\frac{1}{2}$ Kondo system (where the Kondo resonance splits in two peaks, separated by twice the Zeeman energy\cite{2}) we find a fundamentally different splitting. At $B = 1.5 \text{T}$ (Fig. 3c), multiple split peaks appear in regions I and III as ridges of enhanced $dI/dV$ parallel to the $V_T$ axis. In region I, the large zero-bias resonance opens up in four peaks that move linearly with $B$ and become progressively smaller (Fig. 3d). The two inner peaks are due to Zeeman splitting, that is, to higher-order co-tunnelling from $|\downarrow\rangle$ to $|\uparrow\rangle$, $|\uparrow\rangle$ to $|\downarrow\rangle$ (the lower-energy orbital). The two outer peaks arise from co-tunnelling from orbital $|\uparrow\rangle$ to orbital $|\downarrow\rangle$. In the latter case, inter-orbital co-tunnelling processes can occur either with or without spin flip. (However, the corresponding substructure\cite{3} is not resolved owing to the broadening of the outer peaks.) Similar multiple splittings of the Kondo resonance have also been observed in several other samples. According to recent calculations\cite{4}, the observed multiple splitting of the Kondo resonance constitutes direct evidence of SU(4) symmetry, which implies the concomitant presence of spin as well as orbital Kondo correlations, confirming our previous finding.

The slope $|dV/dB|$ of a conductance peak (Fig. 3d) directly yields the value of the magnetic moment associated with the splitting.

We obtain a spin magnetic moment $\mu_{\text{spin}} = |dI/dV|_{|I|} / |I|$ from the inner peaks, and an orbital magnetic moment $\mu_{\text{orb}} = |dI/dV|_{|I|} / \cos \phi = 0.6 \text{meV} \text{T}^{-1} \sim 13 \mu_B$ from the outer peaks ($\phi$ is the angle between the nanotube and $B$)\cite{5}. The same value of $\mu_{\text{orb}}$ follows from the splitting of the Kondo resonance in region III (Fig. 3c). In this case, however, no Zeeman splitting is observed. Here, the magnetic field induces a transition from SU(4) to a spin-based SU(2) Kondo effect for which $k_B T_K$ remains larger than the Zeeman energy. Although the splitting of the Kondo resonance up to a few teslas. Finally, we note that both the one-electron SU(4) and the two-electron singlet–triplet Kondo effects are characterized by a fourfold degeneracy, which results in an enhanced $T_K$ (ref. 27). Apart from this, the two phenomena are fundamentally very different. The singlet–triplet Kondo effect is a spin phenomenon in which the role of the orbital degree of freedom is simply to provide the basis for the construction of spin-singlet and triplet two-particle states (see also Supplementary Information).

Because orbital Kondo correlations can arise only if the orbital quantum number is conserved during tunnelling, our experimental finding of orbital Kondo physics in CNT QDs raises an interesting question concerning the nature of the dot–lead coupling. In our devices, the metal contacts are deposited on top of the CNT and the QD is situated in the segment between them\cite{6}. It is possible that when electrons tunnel out of the QD, they enter first the nanotube section underneath the contacts, where they dwell for some time before moving into the metal. Because the orbital quantum number is probably conserved in a CNT–CNT tunnel process, this intermediate step may account for the observed orbital Kondo effect.

Received 22 November 2004; accepted 25 January 2005; doi:10.1038/nature03422.

Compactly packed photonic crystal fibres

Compact, stable and efficient all-fibre gas cells using hollow-core photonic crystal fibres

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Gas-phase materials are used in a variety of laser-based applications—for example, in high-precision frequency measurement \(^{1,2}\), quantum optics and nonlinear optics \(^{3,4}\). Their full potential has however not been realized because of the lack of a suitable technology for creating gas cells that can guide light over long lengths in a single transverse mode while still offering a high level of integration in a practical and compact set-up or device. As a result, solid-phase materials are still often favoured, even when their performance compares unfavourably with gas-phase systems. Here we report the development of all-fibre gas cells that meet these challenges. Our structures are based on gas-filled hollow-core photonic crystal fibres, in which we have recently demonstrated substantially enhanced stimulated Raman scattering \(^{5,6}\), and which exhibit high performance, excellent long-term pressure stability and ease of use. To illustrate the practical potential of these structures, we report two different devices: a hydrogen-filled cell for efficient generation of rotational Raman scattering using only quasi-continuous-wave laser pulses; and acetylene-filled cells, which we use for absolute frequency-locking of diode lasers with very high signal-to-noise ratios. The stable performance of these compact gas-phase devices could permit, for example, gas-phase laser devices incorporated in a ‘credit card’ or even in a laser pointer.

Our all-fibre gas cells consist of hollow-core photonic crystal fibre (HC-PCF) filled with gas and spliced hermetically at both ends to standard single-mode optical fibre (SMF) (Fig. 1a, d). These novel devices have no bulk-optics components, and may be used with a wide range of commercially available optical fibre components (such as couplers, filters, mirrors and lasers). Moreover, the use of HC-PCF in which light is guided in a single transverse mode by means of a photonic bandgap created in the ‘photonic crystal’ cladding greatly increases the efficiency of these laser–gas devices \(^{3,6}\).

The high air-filling fraction of the HC-PCFs (which are formed by a network of glass webs typically just a few hundred nanometres thick) makes it very challenging to fusion-splice these fibres without collapsing and deforming the microstructure, which would give a high insertion loss. It is also important to avoid contaminating the PCF with solid deposits and water. Nevertheless, after some practice, routine splicing of HC-PCF to SMF was achieved with a typical loss of 1–2 dB. We estimate that a perfect splicing procedure using our fibres would yield a loss of 0.6–0.8 dB. Of this, 0.15 dB arises from the refractive index mismatch between the fibre cores, and is therefore fundamental. The rest is due to modal field mismatch, which we estimate by the butt-joint approximation \(^{7}\) to be 0.4–0.6 dB. The discrepancy between estimated and achieved splice losses is linked to the formation of a recess in the end face of the HC-PCF when heated in the splicer (Fig. 1b). We believe this results from the action of surface tension along the many glass–air interfaces within the holey structure, the viscosity of which offers much less resistance to deformation than the solid glass in the outer cladding of the fibre. The mechanical strength of the splices is...