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Non-interacting 1D conductors

In a one-dimensional (1D) conductor, electrons are confined to move along a single direction, occupying only the quantum mechanical ground state orbital of the transverse dimensions of the wire. What is the electrical conductance of such a quantum wire? This fundamental question was answered by Rolf Landauer many years ago for non-interacting electrons in a clean, ballistic conductor: each spin species carries the quantum of conductance, e^2/h [1], with e the electron charge and h the Planck constant. For a spin degenerate 1D conductor with a single subband, the conductance is therefore $2e^2/h$. If the spin degeneracy is broken and transport of one spin direction is blocked, the conductance is thus reduced to $1e^2/h$. Similar to spin, other degeneracies such as valley degeneracies or multiple 1D subbands due to weaker confinement can also open additional conductance channels.

Conductance quantization is thus a hallmark effect of ballistic 1D noninteracting electrons and was first experimentally demonstrated in gate-defined quantum point contacts in a GaAs 2D electron gas in 1988 [2, 3]. The conductance increases in steps of $2e^2/h$ upon changing the width of the constriction with gate voltage, corresponding to population of 1D-subbands. Theoretically, it can be described in the framework of the Landauer-Büttiker formalism [1, 4]. The conductance quantization is independent of material and sample details, depending only on the number of subbands and the degeneracies present – thus referred to as *universal* conductance quantization. It is also closely related to the quantum Hall effect, where in a strong magnetic field, 1D modes appear at the edge of the sample, each carrying a quantum of conductance.

Interacting 1D conductors: Luttinger liquids

Electrons confined in 1D are genuinely different from free, non-interacting particles. This becomes relevant when replacing the short constrictions ("point contacts") of the first experiments with long wires, where the electrons are tightly confined within a single transverse mode for many Fermi wavelengths along the wire. Due to the Pauli principle, electrons cannot freely pass from one side of the conductor to the other. Instead, they immediately collide with their neighbors and due to the strong 1D confinement, they cannot pass around them. Thus, the effect of disorder and electron-electron (e-e) interactions is very much enhanced in 1D compared to higher dimensions. As a consequence, in clean wires, the electron motion is characterized by density waves arising from the collisions between neighboring electrons, and interactions between them lead to a strong renormalization of the properties of these collective, strongly correlated modes. What emerges is a Luttinger liquid (LL)

[5–8] exhibiting remarkably different physics compared to the Fermi liquid (FL) physics characteristic of 2D and 3D conductors.
Salient signatures of LL theory include ubiquitous power-law scaling, separation of spin and charge modes, and charge fractionalization, all recently observed in experiments [9, 10] performed on cleaved edge overgrowth (CEO) GaAs quantum wires [11, 12] (see box on cleaved edge overgrowth wires). CEO wires are one of the best realizations of a LL liquid known in nature. How do the e-e interactions and the resulting LL physics affect the conductance of the wire? For a clean LL of infinite length, the conductance is renormalized to $K_C\, 2e^2/h$ for a spin degenerate system [13, 14], in principle allowing extraction of the LL interaction parameter K_C in the charge sector from the conductance. The corresponding interaction parameter K_S in the spin sector is normally fixed at $K_S = 1$ due to the absence of significant spin-spin interactions. For repulsive e-e interactions, $0 \leq K_C \leq 1$, where $K_C = 1$ corresponds to non-interacting electrons, $K_C < 0.5$ for long range interactions, and $K_C \to 0$ for interactions approaching infinite strength. The velocity of the LL charge modes is increased to v_F / K_C above the bare Fermi velocity v_F due to increasing stiffness in presence of repulsive e-e interactions $K_C < 1$.

Experiments?

In any realistic experiment, the length of the LL will of course be finite, and FL leads will be attached to the LL at some point. Surprisingly and remarkably, for this more realistic scenario, theory predicts that the *universal* conductance quantization $2e^2/h$ is recovered for the clean, ballistic LL [15–18], irrespective of the strength of interactions K_C. This result can be understood in simple terms when considering that the resistance of a clean 1D quantum wire is really a *contact* resistance arising from the coupling of the higher dimensional modes into the single 1D wire mode, leading to back scattering and thus resistance. This contact process is occurring entirely outside the LL and thus contains no information about the LL and K_C. If the wire itself is ballistic, no further backscattering is caused inside the LL and thus no additional resistance appears. If some even weak disorder is present in the wire, however, the conductance is reduced with typical LL power laws [19, 20].
How do these theoretical predictions compare to experiments? Yacoby and coworkers in Ref. 21 have measured the conductance of ballistic CEO wires already in 1996. They find very well developed plateaus exhibiting quantized conductance as the number of transverse modes is controlled with a gate voltage. To everyone's surprise, however, the conductance quantization was not in units of $2e^2/h$, but rather a lower conductance step was seen, reduced as much as 25% below the universal values at a temperature $T = 0.3$ K. Interestingly, the reduction depended on temperature and source-drain bias V_{SD}, approaching the universal values at high T or V_{SD}. Subsequent attempts to understand this conductance suppression in terms of poor 2D-1D coupling and other possible explanations were inconsistent with important aspects of the data [21, 22]. Thus, the reported non-universal conductance quantization has remained unexplained and has presented an unresolved mystery ever since.

GaAs CEO wires at ultra-low temperatures

In a recent experiment published in 2014 in Physical Review Letters by C. P. Scheller et al. [23] and performed at the University of Basel in an international collaboration with Harvard University (A. Yacoby and G. Barak) and Princeton University (L. N. Pfeiffer and K. W. West), we have revisited the conductance quantization in very similar GaAs CEO wires (see CEO box), now performing experiments for the first time down to 10 mK. Previously, CEO wires were measured at 300 mK or higher temperatures [9, 10, 21, 22]. Obtaining ultra-low temperatures far below 100 mK is rather difficult and has required a significant experimental effort towards filtering and thermalizing the sample and its electrical wires, see [24] for details. Ultra-low temperatures were demonstrated both in-situ on the CEO samples via thermal activation of fractional quantum Hall states as well as in metallic Coulomb blockade thermometers operated under the same conditions, giving an electron temperature of 10.5 ± 0.5 mK at dilution refrigerator temperature $T = 5$ mK. Advancing to ever lower temperatures in quantum transport experiments in nanoscale samples is an ongoing effort at

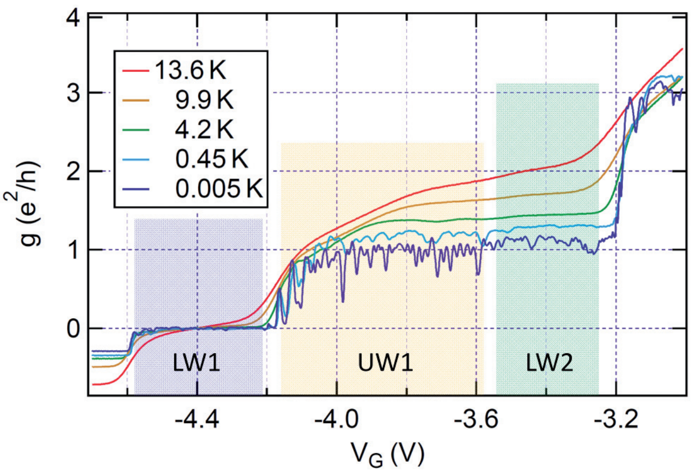


Figure 1: Conductance plateaus of a CEO double wire as a function of temperature

Gate voltage traces of the differential conductance $g(V_G)$ are shown at temperatures T as labeled. Colored bands indicate wire-mode populations: purple: only first mode of the lower wire is occupied (LW1), yellow: first upper wire mode is added (UW1) and green: second LW mode is added (LW2). At the highest T, the thermally smeared remainder of the UW1 plateau approaches $2e^2/h$ (red curve). In addition, at elevated temperatures, we observe a feature reminiscent of the '0.7' structure [26] (shoulder of suppressed g at lower end of plateau). For the lowest T, however, the UW1 conductance plateau is reduced strongly to $1e^2/h$. This is contrary to the expected T-dependence of a 0.7-feature, which rises to $2e^2/h$ at low-T [26].
Conductance oscillations appear at the lowest temperatures, suppressing g below a flat plateau. These are understood as Fabry-Perot type quantum interference arising from scattering outside the ballistic 2 µm wires. Full transmission ~100% is obtained at the maxima of the g-oscillations, which are taken as the relevant measure δg of the wire conductance (see [23] for details). Traces are shifted in g to align LW1 plateaus at g = 0 in order to subtract the LW contribution to the conductance. While UW-LW tunneling is very weak in the 2 µm short gated segments, much stronger UW-LW coupling results for the semi-infinite UW-LW overlap next to the gated region. This LW contribution is subtracted in order to obtain the conductance δg of the UW first mode only, which is adiabatically coupled to the reservoirs. Note that the LW g-contribution is independent of gate voltage V_G since the very long UW-LW overlap is not under the gate.

the University Basel, ultimately striving for the microkelvin temperature range by incorporating advanced nuclear refrigeration schemes [25], with the goal of opening the doors for new physics.

The results reported in [23] are rather striking: The conductance of the first wire mode reaches $1e^2/h$ at $T \sim 100$ mK and remains fixed at this value for lower T, while the electron temperature falls far below 100 mK. At high $T \gtrsim 10$ K, the conductance approaches the expected universal value $2e^2/h$ (see Fig.1 and 2). This suggests lifting of the electron spin degeneracy, yet without an external magnetic field. The same behavior was seen in all investigated samples, is robust against variation of the wire electron density and persists at moderate magnetic fields (up to 3 T). Further, application of even a small source-drain bias voltage acts to destroy the low-conductance state, driving the conductance back towards $2e^2/h$, similar to temperature. This suggests the emergence of a new, small energy scale in the physics of the wire. We emphasize that in the high temperature range $T \geq 0.3$ K, our results are fully consistent with the previous experiment [21].

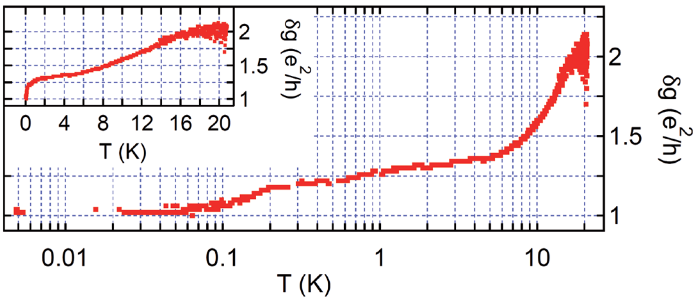


Figure 2: Conductance reduction by a factor of 2

Conductance δg of the first mode of the UW as a function of temperature on a logarithmic axis (linear axis in inset), extracted from conductance plateau traces as in Fig. 1. Small but discrete steps in g result from a histogram binning effect. We note that the transition from $2e^2/h$ to $1e^2/h$ occurs over a very broad temperature range spanning about 2 orders of magnitude in temperatures. Such broad cross-overs are a characteristic signature of LL physics.

A detailed analysis of the data has been performed [23], considering all possible models we were aware of, including non-interacting electrons, e-e interactions only within the wire and also in the 1D electron gas outside the wire (variations of LL physics), poor 2D-1D coupling in presence of LL correlations, an incoherent LL due to Wigner crystal formation and, finally, also the effects of spin-orbit coupling. While some of these models can capture certain aspects of the data, clear and significant inconsistencies appear with salient features of the experiment. Ultimately, all of these models had to be rejected. Only a recent theory of helical nuclear spin order in a LL by Braunecker, Simon and Loss [27] can account for the experimental observations without inconsistencies, remaining as the figurative "last theory standing".

Helical Nuclear Spin Order

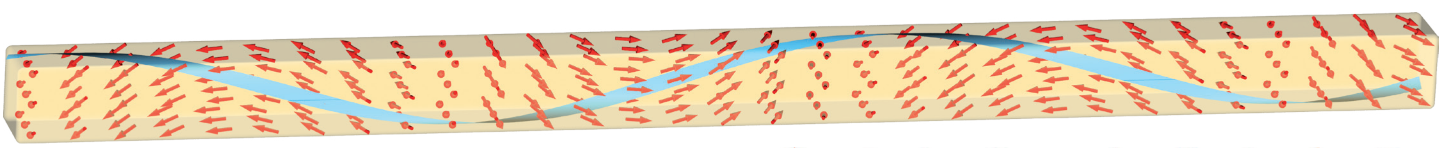
The conceptual advance made by this recent theory [27] may be daring and profound, yet it is at the same time simple and natural: our discussion of 1D conductors and GaAs wires in particular so far has neglected the nuclear spin. In fact, all stable isotopes of both Ga and As have nuclear spin $I = 3/2$. Each transverse wire cross section contains 10^3 to 10^4 nuclear spins which can couple to the electron spin via the hyperfine interaction, defining a new type of central spin problem in a LL. Thus, in this sense, the wires investigated in the experiments operate in a "quasi"-1D regime where a single electronic mode couples to a large number of nuclear spins. The theory then proceeds to calculate the consequences of this, predicting profound and non-trivial results, outlined here.

Below a cross-over temperature T^*, an effective Ruderman-Kittel-Kasuya-Yosida (RKKY) interaction, strongly enhanced by e-e interactions of the 1D electronic modes, forces the nuclear spin system via the hyperfine interaction into helical order [27] (see also box 2). The resulting large Overhauser field with spiral texture in space acts back on the electronic system where a large gap opens – pinned at the Fermi energy – for half of the low energy electronic modes. A helical (spin-filtered) LL thus forms and causes the reduction of the conductance by a factor of 2 in the absence of an external magnetic field, applicable similarly for single and double wires [28] and even for arrays of wires [29]. A novel state of strongly correlated quasi-1D quantum matter is therefore predicted, where nuclear spins and the half-gapped electron spins are locked into order in a conjoined helical spin state of perfectly synchronized and phase locked spiraling of electron and nuclear spins, as depicted in Fig. 3. This new state is a thermodynamic ground state of the system protected by a gap, rather than a dynamic polarization effect caused by driving out of equilibrium. Finally, the spin helix state is a clear manifestation of LL physics present in the electronic sector. In absence of LL correlations, the ordering temperature quickly drops to much lower temperatures.

Comparison to Experiment

The cross-over temperature T^* predicted by theory depends strongly on the LL interaction parameter K_c. While neither determining K_c from experiment nor estimating it theoretically is trivial, reasonable values [9, 10] for the single mode case are approximately 0.4, resulting in a $T^* \sim 0.2$ K, while $K_c = 0.3$ already gives $T^* \sim 0.6$ K. Considering that full nuclear order and $g = 1e^2/h$ is obtained only at $T \ll T^*$ and zero polarization with $g = 2e^2/h$ only at $T \gg T^*$, these ordering temperatures are consistent with the experiment.

Further, a very broad, washed out transition occurring over a large range of temperatures would be expected for a LL system – as observed here in the experiment. In addition, a Zeeman splitting smaller than the induced electronic gap



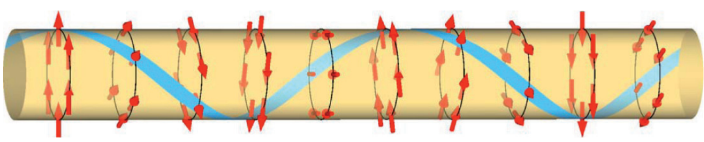


Figure 3: Schematic representation of helical electron spin order (blue) and nuclear spin order (red arrows) in a GaAs quantum wire (upper panel) and alternatively in a ^{13}C carbon nanotube (nuclear spin 1/2, lower panel). Lower panel taken from [27].

should affect neither the nuclear order nor the conductance strongly, consistent with the observed insensitivity to moderate magnetic fields (up to 3 T). Finally, strong sensitivity to source-drain bias could be related to the energy to destroy the nuclear spin helix. However, this could also have other origins, including resistive heating.
Several clear characteristics are present in the data that support the nuclear spin helix model without contradiction: the conductance reduction by a factor of 2 followed by a saturation, a crossover temperature in the observed range, a very broad transition, sensitivity to source-drain bias and finally, insensitivity to both a small Zeeman splitting and change of density. Nevertheless, all present data stem from electronic transport measurements, and no direct evidence for nuclear spin order is currently available. Further experiments are required to learn more about this system.

Future Experiments?

Directly probing the nuclear spin helix with a scanning magnetometer would be very difficult: the wires are buried hundreds of nanometers below the surface, the nuclear magnetic moments are tiny (with only few electrons present in the wire), and there is no overall magnetization to be detected due to the spiraling and self-canceling nature of the polarization. Further, we note that it is difficult to estimate what the effect of an NMR type excitation is on the system, what the low energy nuclear spin excitations are, and whether a detectable resistive signal would result. Work is currently under way to investigate these and other questions, both in theory and experiment. Tunneling spectroscopy with two parallel wires can be used to map the electronic dispersions of the wires [9], and might be used to observe the partial spin gap below T^*. In any case, the data presented in [23] are striking and stand alone, irrespective of the model used for interpretation.

Hottest nuclear spin order

If nuclear spin order was indeed observed in the experiment [23], then this would constitute by far the highest temperature at which nuclear order was reported to date. Due to the tiny size of the nuclear magnetic moment, nuclear dipole-dipole interactions are extremely weak, leading to ordering only at extremely low temperatures, typically at microkelvin or lower temperatures. However, also in 3D bulk systems, enhancement of nuclear spin-spin interactions via hyperfine coupling and conduction electron RKKY mechanism have

GaAs cleaved edge overgrowth (CEO) quantum wires

CEO wires are fabricated starting from a high-mobility 2D electron gas (2DEG) in a GaAs quantum well by cleaving the wafer inside the ultra-high vacuum MBE chamber and over-growing another modulation doping sequence on the freshly cleaved edge. The additional Si doping from the overgrowth at the edge combined with rearrangement of the resulting band structure leads to accumulation of charges and formation of quantum modes along the CEO edge. Thus, a quantum wire is created running along the 2DEG cleaved edge forming a 1D electron gas (1DEG) consisting of a few transverse modes, see [11, 12] for fabrication details.
Along similar lines, double wire (DW) samples featuring two parallel quantum wires at the cleaved edge can be created. Here, only the upper 2DEG is doped and populated in the double quantum well wafer, while both upper wire (UW) and lower wire (LW) at the edge are populated and conducting (see Fig. 4). The LW is only weakly tunnel coupled to the UW through a 6 nm thick AlGaAs barrier, giving a tunneling conductance of 0.03 e^2/h at zero B-field for a 2 µm long segment. Hence the 2 µm long DWs are considered as independent parallel resistors, with total conductance given by the sum of each conductance.
A top gate of 2 µm length allows for local depletion of the 2DEG below it (see Fig.4), thus creating wire segments isolated from the 2DEG under the gate, as desired. Further, the gate voltage can control the number of modes in both the upper and lower wire under the gate (though not separately in each), down to the single mode LL regime. The gated 2 µm wire segments – in the following referred to as the "wires" – are extending into the ungated, semi-infinite 1DEGs on each side (samples are about 5 mm long), which in turn are connected to the adjacent 2DEG and its ohmic contacts, thus allowing for measurements of the wire conductance. The 2DEG-1DEG contact of effectively semi-infinite size together with a gate-defined and smooth *1DEG* to single-mode wire transition assure an overall adiabatic coupling to the LL wire.
The resulting quantum wires are of exceptional quality, probably the best realization known today of a clean 1D conductor. Beyond gate-control of the charge density and transverse modes populated, a mean free path of more than 10 µm and a very large transverse subband spacing of ~ 20 meV make these wires ideal for studying e-e interactions in 1D and LL physics. Such wires display the probably strongest evidence for LL physics to date, such as spin-charge separation [9] and charge fractionalization [10]. However, CEO samples are extremely difficult to produce and prepare for measurement, forcing our present experiment to resort to the limited stock of wires currently available. Due to the lower quality of single wire samples (fabricated more than 15 years ago), we have focused our measurements primarily on more recently fabricated clean double wires. (The Wegscheider group (ETH Zürich) is recently working on fabricating new CEO wires.)

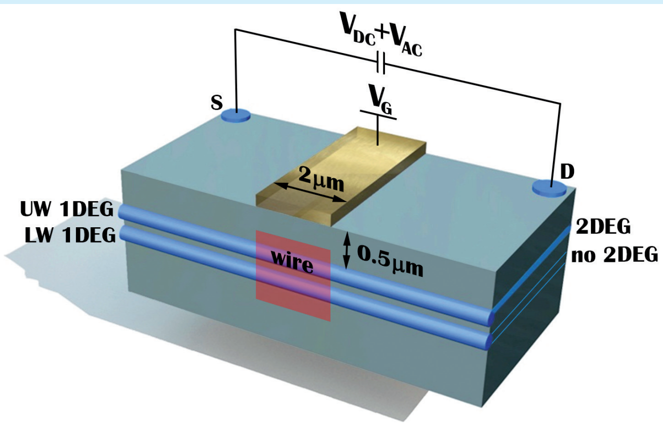


Figure 4: Double-wire device schematic.

been observed [30], for example in metals or strongly correlated conductors, giving ordering temperatures as high as 0.4 mK (2.6 mK) for the hyperfine-enhanced rare-earth Van-Vleck compounds $PrNi_5$ ($PrCu_6$), respectively [31]. The mechanism is similar to the one described above. However, in a 1D conductor (LL), the e-e interactions are very much enhanced compared to 3D. Further, strong and intricate feedback between electronic and nuclear modes further enhances the ordering temperature. Together, this leads to spontaneous nuclear order already in the 100 mK range for the GaAs LL wire, at two orders of magnitude higher temperatures compared to bulk systems.

Another interesting aspect of the GaAs wire is its capability to be gated: the electrons can be depleted from the wire with

Electron and nuclear spin helices: Peierls instability and RKKY interactions

The combined ordered state of nuclear spins and electrons arises through a feedback process closely connected with the Peierls instability, a generic effect in any 1D conductor. The Peierls instability is best explained by considering the upper part of Fig. 5 representing the band structure of a 1D conductor. If the conductor is exposed to a periodic potential with spatial period of half the Fermi wavelength, electron scattering on this potential leads to a momentum transfer of $2k_F$, with k_F the Fermi momentum. The electron states near both Fermi points at $\pm k_F$ mix due to the periodic potential, a gap opens, and the system becomes insulating. Exposing the electron system to a spiral magnetic field like potential with period of half the Fermi wavelength, as created by a nuclear spin helix, leads to a *spin-selective Peierls transition* [31]. Scattering on the helix has the same effect of inducing a mixing of the states at $\pm k_F$, yet with the additional effect that a momentum transfer of $+2k_F$ is accompanied by an upflip of the electron spin, and a $-2k_F$ transfer with a downflip (or vice versa if the helix rotates in the opposite direction). Consequently, only half of the electron modes can undergo the Peierls transition and become insulating, while the other half with the non-matching spins remains conducting and forms a helical (spin-filtered) conductor.

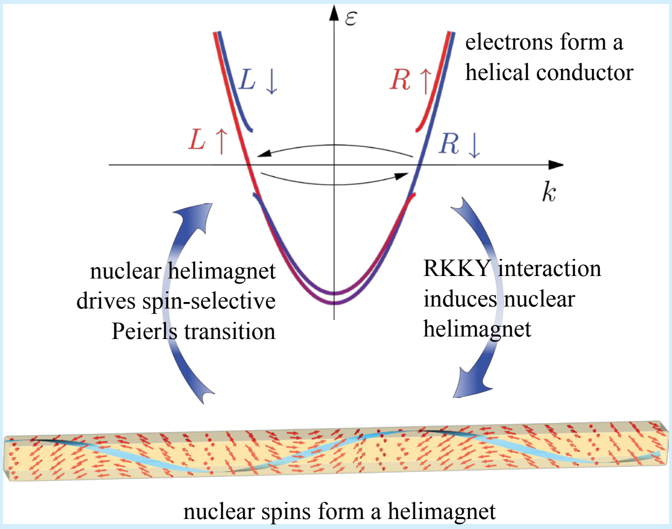


Figure 5: Illustration of the feedback mechanism stabilizing the joint nuclear spin and electron order. The RKKY interaction between nuclear spins mediated by the electron conductor induces nuclear magnetic order in form of a spiral with spatial period of half the Fermi wavelength. The backaction of the periodic magnetic potential formed by the nuclear spins causes selective spin-flip scattering between the two Fermi points, opening a gap for one half of the opposite spins and leaving the other half in a spin-filtered, helical conducting state. In turn, the RKKY interaction becomes stronger and the joint ordered state is strongly stabilized by this feedback effect, together with a strong renormalization by e-e interactions.

The Peierls instability is also the origin of the nuclear spin helix. The dominant interaction between the nuclear spins is the Ruderman-Kittel-Kasuya-Yosida (RKKY) interaction, which is an indirect long-range interaction mediated by the electron system. A spin exchange between a nuclear spin and an electron spin creates locally a magnetic excitation in the electron system that propagates through the conductor and can induce a further spin exchange with a distant nuclear spin. The result is an effective Heisenberg interaction between the nuclear spins with an interaction strength given by the electron response function for magnetic excitations, the electron spin susceptibility. In momentum space, the latter is strongly peaked at $2k_F$ (for normal as well as for helical conductors), which is nothing but the manifestation of the Peierls instability, and it shows that the energetically most favorable state for the nuclear spins is to form a $2k_F$ spatial modulation. A detailed analysis of the ordering transition and its stability then reveals that the nuclear spins order in form of the helix shown in Figure 5, with clockwise or anticlockwise helicity as well as the plane in which the spins rotate selected by a spontaneous symmetry breaking. The spin-selective Peierls transition enhances the strength of the RKKY interaction, and this feedback between the subsystems, together with a strong renormalization of the coupling strengths by electron interactions, leads to a strong stabilization of the combined ordered state of nuclear spins and electrons.

Traditionally, order between electron and spin systems can be split into two classes. First the class, in which the spins and the electrons form a joint strongly correlated state, such as in Kondo lattice systems at temperatures below the Kondo temperature [32]. Second the class of the type of nuclear magnets in three-dimensional metals [33], in which the nuclear spins order due to the presence of the electrons, yet the electrons themselves are unaffected by the nuclear spins. The state shown in Fig. 5 forms an intermediate class, in which nuclear spins and electrons order individually but are tightly bound together through a self-consistent feedback mechanism [27]. In this case, the nuclear spins order due to their effective RKKY interaction, which is mediated through the electrons, but do not form a coherent correlated state with them. Yet through their ordering they generate a magnetic superstructure, the nuclear spin helix, that acts back on the electrons. This back action triggers the formation of a spiral electron spin density wave for half of the conduction electron modes. The other half remains conducting in a strongly renormalized helical (spin-filtered) state and further stabilizes the nuclear helix.

a gate voltage, hence also removing the nuclear order. One can thus gain electrical control over nuclear spin order.

Future prospects of helical Luttinger liquids

The prospects of helical nuclear spin order are exciting and far reaching for a number of reasons: first, the ordered wire is a helical conductor with opposite spins moving in opposite directions, therefore acting as an excellent spin filter. Second, in the ordered phase, any nuclear spin fluctuations would be fully suppressed, thus eliminating the predominant source of decoherence in GaAs electron spin qubits (if such a qubit could be realized in the 1D wire). This was predicted earlier [27] (see also the inaugural article of the *Progress in Physics* series, from April 2007, *http://www.sps.ch/artikel/progresses/the_pleasures_on_the_road_to_a_quantum_computer_1/*) and served as one of the original motivations for work on this subject.

Last but not least, a helical LL brought into contact with a BCS superconductor can serve as a platform for Majorana fermions exhibiting non-Abelian braiding and allowing for fault-tolerant topological quantum computing. It has been shown [31] that the helical LL induced here by a nuclear spin helix is equivalent to a Rashba spin-orbit (SO) wire in presence of a magnetic field which opens a $k = 0$ helical gap. The strength of the equivalent Rashba-type SO interaction is given essentially by the Fermi wave number k_F, i.e. corresponds to a spin-orbit length of about $2\pi / k_F \lesssim 100$ nm – a very strong and thus useful SO interaction. Unlike the Rashba wire, the gap of the helical LL is always pinned at the Fermi energy. This is a clear advantage over the Rashba wire, where the chemical potential (density in the wire) has to be carefully tuned into the gap (fixed at an energy defined by the wire SO coupling strength).
While inducing a superconducting proximity effect in a GaAs CEO wire is a challenging, maybe daunting goal, it is maybe conceivable to add a thin Al layer (superconducting critial temperature of 1.2 K) to the overgrowth sequence – Al is always needed and available in high-mobility GaAs MBE systems. With magnetic field applied in the plane of the Al layer, the corresponding critical field can exceed the bulk critical field of ~ 10 mT by orders of magnitude, thus potentially allowing for the creation of a topological phase in a magnetic field, with Majorana fermions at its ends. We hope that the experimental results already presented here are an important step in the direction of new developments in this exciting field of research.

References

[1] R. Landauer, IBM J. Res. Dev **1**, 223 (1957).
[2] B. J. van Wees, H. van Houten, C. W. J. Beenakker, J. G. Williamson, L. P. Kouwenhoven, D. van der Marel, and C. T. Foxon, Phys. Rev. Lett. **60**, 848 (1988).
[3] D. A. Wharam, T. J. Thornton, R. Newbury, M. Pepper, H. Ahmed, J. E. F. Frost, D. G. Hasko, D. C. Peacock, D. A. Ritchie, and G. A. C. Jones, J. Phys. C **21**, L209 (1988).
[4] M. Büttiker, Y. Imry, R. Landauer, and S. Pinhas, Phys. Rev. B **31**, 6207 (1985).
[5] S. I. Tomonaga, Prog. Theor. Phys. **5**, 544 (1950).
[6] J. M. Luttinger, J. Math. Phys.(N.Y.) **4**, 1154 (1963).
[7] F. D. M. Haldane, J. Phys. C **14**, 2585 (1981).
[8] T. Giamarchi, *Quantum Physics in One Dimension* (Oxford University Press, 2004).
[9] O. M. Auslaender, H. Steinberg, A. Yacoby, Y. Tserkovnyak, B. Halperin, K. Baldwin, L. Pfeiffer, and K. W. West, Science **308**, 88 (2005).
[10] H. Steinberg, G. Barak, A. Yacoby, L. Pfeiffer, K. West, B. Halperin, and K. L. Hur, Nat. Phys. **4**, 116 (2008).
[11] L. N. Pfeiffer, H. L. Stormer, K. W. Baldwin, K. W. West, A. R. Goñi, A. Pinczuk, R. C. Ashoori, M. M. Dignam, and W. Wegscheider, J. Cryst. Growth. **127**, 849 (1993).
[12] W. Wegscheider, W. Kang, L. N. Pfeiffer, K. W. West, H. L. Stormer, and K. W. Baldwin, Solid-State Electronics **37**, 547 (1994).
[13] W. Apel and T. M. Rice, Phys. Rev. B **26**, 7063 (1982).
[14] C. L. Kane and M. P. A. Fischer, Phys. Rev. Lett. **68**, 665 (1992).
[15] D. L. Maslov and M. Stone, Phys. Rev. B **52**, R5539 (1995).
[16] I. Safi and H. J. Schulz, Phys. Rev. B **52**, R17040 (1995).
[17] V. V. Ponomarenko, Phys. Rev. B **52**, R8666 (1995).
[18] Y. Oreg and A. M. Finkel'stein, Phys. Rev. B **54**, 14265 (1996).
[19] S. Tarucha, T. Honda, and T. Saku, Solid State Commun. **94**, 413 (1994).
[20] D. L. Maslov, Phys. Rev. B **52**, R14368 (1995).
[21] A. Yacoby, H. L. Stormer, N. S. Wingreen, L. N. Pfeiffer, K. W. Baldwin, and K. W. West, Phys. Rev. Lett. **77**, 4612 (1996).
[22] R. de Picciotto, H. L. Stormer, A. Yacoby, L. N. Pfeiffer, K. W. Baldwin, and K. W. West, Phys. Rev. Lett. **85**, 1730 (2000).
[23] C. P. Scheller, T. M. Liu, G. Barak, A. Yacoby, L. N. Pfeiffer, K. W. West, and D. M. Zumbühl, Phys. Rev. Lett. **112**, 066801 (2014).
[24] C. P. Scheller, S. Heizmann, K. Bedner, D. Giss, M. Meschke, D. M. Zumbühl, J. D. Zimmerman, and A. C. Gossard, Appl. Phys. Lett. **104**, 211106 (2014).
[25] L. Casparis, M. Meschke, D. Maradan, A. C. Clark, C. P. Scheller, K. K. Schwarzwälder, J. P. Pekola, and D. M. Zumbühl, Rev. Sci. Instum. **83**, 083903 (2012).
[26] K. J. Thomas, J. T. Nicholls, M. Y. Simmons, M. Pepper, D. R. Mace, and D. A. Ritchie, Phys. Rev. Lett. **77**, 135 (1996).
[27] B. Braunecker, P. Simon, and D. Loss, Phys. Rev. B **80**, 165119 (2009); P. Simon and D. Loss, Phys. Rev. Lett. **98**, 156401 (2007).
[28] T. Meng and D. Loss, Phys. Rev. B **87**, 235427 (2013).
[29] T. Meng, P. Stano, J. Klinovaja, and D. Loss, Eur. J. Phys. B, to be published.
[30] A. S. Oja and O. V. Lounasmaa, Rev. Mod. Phys. **69**, 1 (1997).
[31] B. Braunecker, G. I. Japaridze, J. Klinovaja, and D. Loss, Phys. Rev. B **82**, 045127 (2010).
[32] H. Tsunetsugu, M. Sigrist, and K. Ueda, Rev. Mod. Phys. **69**, 809 (1997).
[33] H. Fröhlich and F. R. N. Nabarro, Proc. R. Soc. Lond. A **175**, 382 (1940).