Direction-resolved transport and possible many-body effects in one-dimensional thermopower

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A single-particle theory due to Mott predicts a proportionality between the diffusion thermopower and the energy derivative of the logarithm of the conductance. Measurements of a ballistic 1D wire show that the Mott theory remains valid in the presence of a finite current, and that it leads to a direction-sensitive probe of electron transport. We observe an apparent violation of the Mott model at low electron densities, when there is a nonquantized plateau in the conductance at $0.7(2e^2/h)$. There is as yet no successful theoretical explanation of this so called *0.7 structure*, but the distinctive thermopower signature, which deviates from single-particle predictions, may provide the key to a better understanding.

In linear response, the diffusion thermopower of a noninteracting, degenerate electron gas $S=(\delta V/\delta T)_{I=0}$ is related to the conductance $G=(\delta I/\delta V)_{\delta T=0}$ through the relation[1]

$$S=-\frac{\pi^2 k_B^2 T}{3e}\frac{1}{G}\frac{\partial G}{\partial \mu}, \tag{1}$$

where μ is the chemical potential. Equation (1) remains essentially unchanged when the theory is reformulated for a mesoscopic device connected to Fermi function reservoirs.[2–4] Measurements of S have recently been used to investigate mesoscopic systems, for example, open quantum dots show[5] fluctuations in S with a non-Gaussian spectrum characteristic of chaotic behavior, yielding information which is not accessible from G. We recently[6] confirmed Eq. (1) for ballistic one-dimensional (1D) transport through a GaAs quantum wire, showing that an electron gas can be used as its own thermometer. In this paper we present thermopower results in two new regimes of 1D transport: (i) when a strong electric field lifts the momentum degeneracy, and (ii), when the 1D electron gas is strongly interacting.

Equation (1) assumes that electrons, supplied from reservoirs with Fermi statistics, pass through the 1D wire with no external anisotropy such as phonon drag from a heat flux. Transport is elastic, so that all of the electron's thermal energy is transmitted together with its charge, and no heat originates or is dissipated in the device. This can be achieved in low-dimensional devices, as there is only weak electron-phonon coupling at low temperatures.[6,7] Under these conditions, the net current passing through the device is given by

$$I=-\frac{1}{e}\int_0^\infty dE\, g(E)[f(E|\mu_1,T_1)-f(E|\mu_2,T_2)], \tag{2}$$

where $f(E|\mu_i,T_i)$ is the Fermi function $(1+\exp[(E-\mu_i)/k_BT_i])^{-1}$, and the subscripts $i=1,2$ refer to the two reservoirs. The effective conductance is $g(E)=e^2\mathcal{D}(E)v(E)\tau(E)$, where $\mathcal{D}(E)$ is the density of states, $\tau(E)$ is the electron transmission probability, and $v(E)$ is the group velocity. In 1D the product $\mathcal{D}(E)v(E)=1/h$, so that $g=e^2\tau(E)/h$, which is then summed over all subbands and spins. An expansion of Eq. (2) for small chemical potential and temperature differences between the reservoirs leads to $G=g(\mu)$, and S as given in Eq. (1).

The conductance G of a clean 1D wire shows steps between quantized values as successive 1D subbands are opened to electron transport; the corresponding thermopower S shows a series of peaks. A comparison of the height and shape of these peaks to the conductance characteristics confirms the prediction of Eq. (1). This allows us to deduce[6] the temperature difference across the constriction and the electron energy relaxation rate in a two-dimensional electron gas (2DEG), which agrees well with the theory of phonon scattering.[7] Earlier S measurements[8,9] of a 1D constriction could not be applied in this way, as they were either uncalibrated, or were operated beyond the validity of Eq. (1).

Figure 1 shows schematically how S and G of a 1D wire are measured; further details are given elsewhere.[8,6] The 1D constriction is produced by electrostatically squeezing a 2DEG (at a GaAs/AlGaAs interface) using a gate voltage V_g applied to a pair of Schottky gates on the surface. The 2DEG mesa is patterned so that an AC current I_H at frequency f_H heats the electrons in a wide channel[10] on one side of the constriction to a temperature T_2, while the electrons on the other side of the constriction remain at the lattice temperature T_1.

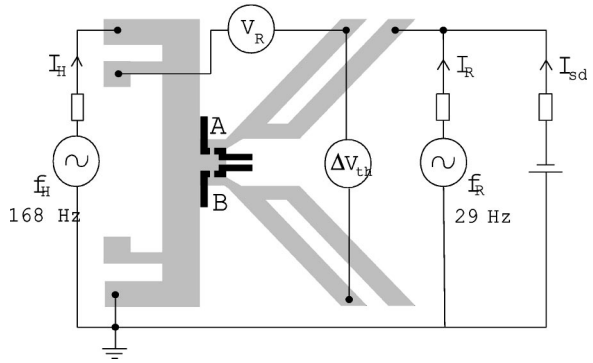

FIG. 1. Simultaneous measurements of thermopower and conductance. An AC current I_H at frequency f_H heats electrons on the left hand side of the sample by ΔT. The difference in the thermopowers of the 1D constrictions A (variable) and B (fixed) leads to a thermoelectric voltage $\Delta V_{\rm th}=(S_A-S_B)\Delta T$ at $2f_H$. The voltage V_R at f_R gives the electrical resistance $R=1/G$ of constriction A, through which a DC current $I_{\rm sd}$ can also be passed.

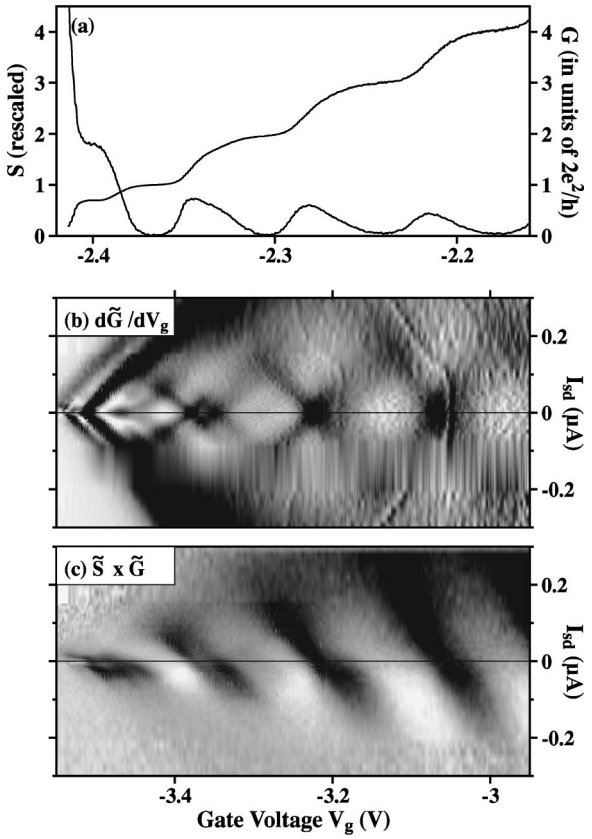

FIG. 2. (a) The conductance $G(V_g)=1/R$ and thermopower S for sample C, showing that S goes to zero on a quantized conductance plateau. The thermopower in this and subsequent figures has been rescaled so that the first peak has unity height. Greyscale plots of DC current-biased measurements of (b) the transconductance $d\tilde{G}/dV_g$, and (c) the product $\tilde{S}\tilde{G}$ for sample A. At $I_{sd}=0$, the two measurements give similar symmetric behavior, but with a finite bias it becomes clear that the thermopower is sensitive only to electron transmission originating from one of the contacts.

High impedance probes and current injection are used, with an electrical earth defined in the heating channel. The heating power varies as I_H^2, so a thermoelectric voltage $\Delta V_{\rm th}$ is generated across the 1D constriction at $2f_H$, allowing thermal effects to be distinguished from resistive voltages. A simultaneous measurement of the conductance of the constriction is made by passing a current at a different frequency (f_R).

Measurements were performed on three samples, each of which had split gates of lithographic length 0.8 μm and gap width 0.8 μm fabricated over a 2700 Å deep 2DEG. After illumination by a red light-emitting diode the electron densities and mobilities were $n=3.6\times10^{11}$ cm^{-2} and $\mu=4.5\times10^6$ cm^2/V s for samples A and C, where the 2DEG is formed in a 200 Å-wide quantum well, and $n=1.8\times10^{11}$ cm^{-2} and $\mu=2.4\times10^6$ cm^2/V s for sample B, which is a single heterojunction. Figure 2(a) shows the measured conductance and thermopower of sample C at $T=300$ mK, as the constriction width is varied using the gate voltage V_g. The conductance G shows plateaus quantized at multiples of $2e^2/h$, between which there are, as expected, peaks in the thermopower S. The thermopower on the first peak is 0.1 μV/K, from which we estimate[6] the temperature difference $T_2-T_1\approx0.1$ K. The measurements are therefore in the linear response regime.

We now consider the effects of a finite DC current. DC voltage biasing of the conductance is an established technique for measuring[11] 1D subband spacings in a ballistic wire. Since G is determined by electrons at the chemical potential of the contacts, a transmission feature will be seen when it aligns with either the source or the drain; compared to equilibrium conductance traces there is a doubling of structure in the presence of a source-drain voltage. In the measurements presented here, we produce a voltage difference, and hence a chemical potential offset, between the source and drain using a DC bias current $I_{\rm sd}$. The linear response to a small additional AC signal gives a differential thermopower $\tilde{S}$

$$\tilde{S}=-\frac{1}{e}\left(\frac{\partial\mu_1}{\partial T_2}\right)_{\mu_2,I=0}=-\frac{\pi^2k_B^2T_2}{3e}\frac{1}{\tilde{G}}\left(\frac{\partial g}{\partial E}\right)_{\mu_2}, \quad (3)$$

and differential conductance $\tilde{G}$

$$\tilde{G}=e\left(\frac{\partial I}{\partial\mu_1}\right)_{\mu_2,(T_1-T_2)=0}=\beta g(\mu_1)+(1-\beta)g(\mu_2), \quad (4)$$

where $\beta\approx1/2$ is the fraction[11] of the AC voltage which is dropped between the 1D constriction and contact 1, and it is assumed that the transmission properties of the constriction are temperature independent.

The measured variation of the transconductance $d\tilde{G}/dV_g$ and the product $\tilde{S}\tilde{G}$ with DC current $I_{\rm sd}$ are shown in the greyscale plots of Figs. 2(b) and 2(c) for sample A. Apart from the interaction effects discussed later, these two quantities are proportional to each other at $I_{\rm sd}=0$. The transconductance, which is high (dark areas) at the transitions between conductance plateaus, shows equal splitting of each peak as the bias is applied. This has been used[11] to calibrate the gate voltage as an energy scale. The second plot, however, shows only one of these branches, as can be understood from Eq. (3): the product $\tilde{S}\tilde{G}$ is determined by the transmission properties of the 1D constriction only at the chemical potential μ_2 of the contact where the AC heating occurs, and is independent of the other contact. In contrast, the two contacts contribute approximately equal weight to the DC biased differential conductance measurement. The differential thermopower $\tilde{S}$ gives direction-resolved information about the electron transport, which is unobtainable from conductance measurements. The data also demonstrate how the Mott formalism can be successfully extended to describe measurements which are far from linear response.

At these low electron densities, just after the last spin-degenerate subband has been depopulated, a nonquantized plateau has been measured[12–14] with a conductance of about $0.7(2e^2/h)$. This *0.7 structure* increases in strength with increasing temperature, suggesting that it arises from an excited state of the 1D electron gas.[14] The plateau value of conductance decreases continuously as a parallel magnetic field $B_\parallel$ is increased,[12] eventually becoming the quantized spin-split plateau at e^2/h. Since the g-factor of the 1D sub-

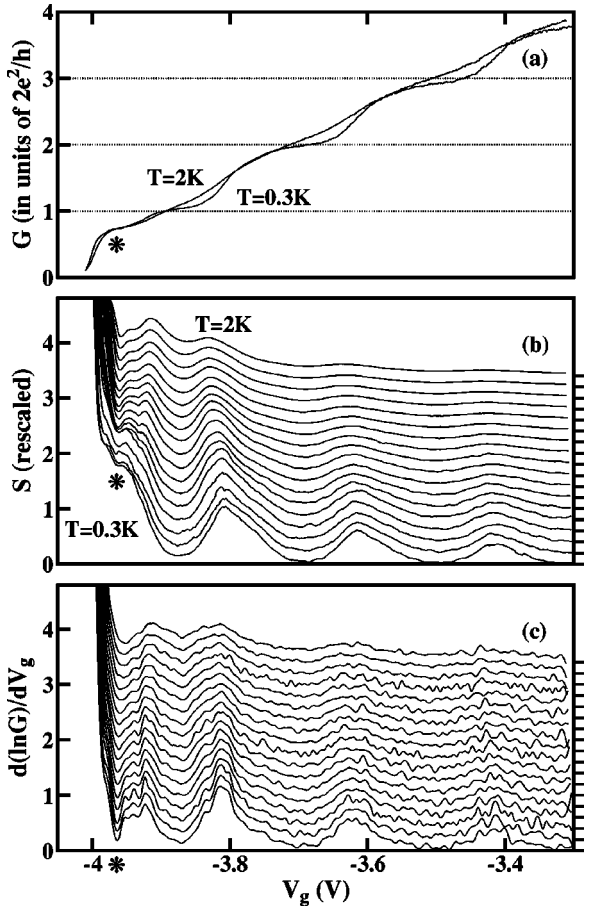

FIG. 3. (a) The conductance G of sample B at $T=0.3$ and 2 K. (b) The thermopower S and (c) the calculated quantity $d(\ln G)/dV_g$ from 0.3 to 2 K in increments of 0.1 K. Agreement is good for a wide channel, but the zero seen in $d(\ln G)/dV_g$ at low temperatures, marked by an asterix *, is absent in S. Traces in (b) and (c) are offset vertically, with zeros as marked on the right-hand vertical axis.

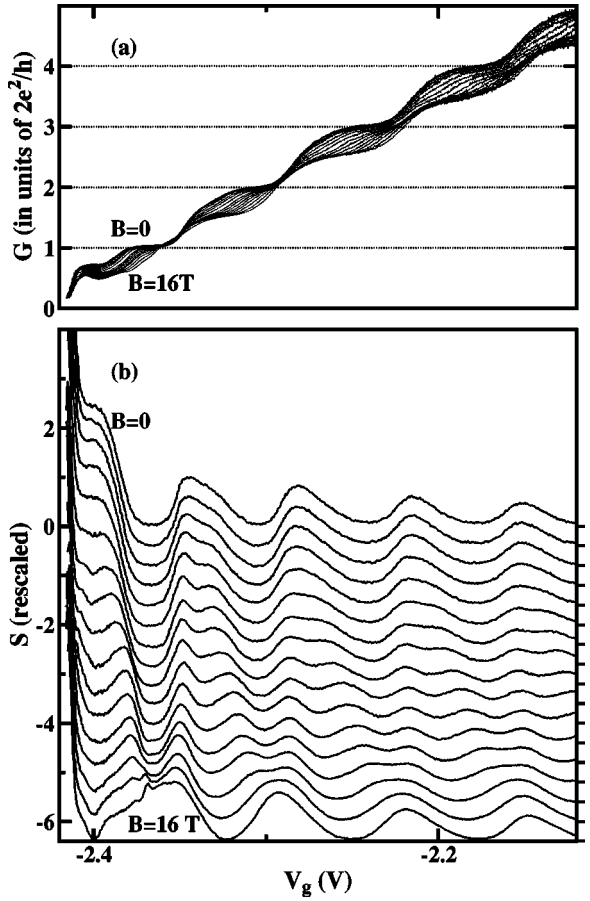

FIG. 4. (a) The conductance G and (b) thermopower S for sample C when the in-plane magnetic field $B_\parallel$ is incremented from 0 to 16 T in steps of 1 T. Lifting of spin degeneracy at high fields restores the zero in S that is predicted by single-particle theory [Eq. (1)]. The traces in (b) are offset vertically.

bands increases[13] as the subband index is reduced, this may point to the importance of many-body spin interactions in determining the behavior. It is worth emphasizing that the 0.7 structure is different from the observed[15] depression of the quantized conductance plateaus in long wires. In the split-gate samples investigated here, the 0.7 structure is measured in addition to the usual plateaus, which remain quantized at multiples of $2e^2/h$.

Figure 3(a) shows the 0.7 structure as a function of temperature in sample B; the corresponding thermopower signal is shown in Fig. 3(b). The prediction of Eq. (1), shown in Fig. 3(c), has been calculated using the gate-voltage derivative of the conductance $\partial G/\partial V_g$ in place of $\partial G/\partial \mu$, assuming that the gate voltage smoothly shifts the 1D confinement energy and provides a linear measure of the mesoscopic energy scale. The calculation, which predicts that a zero in S should accompany a plateau in G, is not in agreement with the measured thermopower. The discrepancy cannot be attributed to thermal broadening, as this would affect both G and S, and in any case there is no change in the shape of the measured thermopower traces below 1 K.

Figures 4(a) and 4(b) show S and G measurements for sample C, as the spin degeneracy is lifted by a parallel magnetic field $B_\parallel$. In high fields, when the spin-split conductance plateau at e^2/h is fully developed, there is a zero in the thermopower S, and the validity of Eq. (1) is restored. The measurement of a finite thermopower when the conductance is on the plateau at $B=0$ indicates a breakdown of the Mott model, and we interpret it as a manifestation of many-body effects in the 1D electron gas. Interactions are most likely to affect the thermopower through a dependence of the transmission, $g(E)$, on the electron density or temperature, which are in turn determined by the properties of the contacts. With this in mind, a rederivation of Eq. (1) suggests that $\partial G/\partial \mu$ should be replaced by $\langle \partial G/\partial E \rangle_{E=\mu}$, where the brackets $\langle\ \rangle$ indicate ensemble averaging over thermal fluctuations.

It is possible that the 1D electron gas within the constriction is incompressible, so that the gate voltage V_g cannot change the position of the subband relative to the chemical potential; this would result in plateaus in G and S, as observed. The measured thermopower will differ from $\partial \ln G/\partial V_g$, because in Eq. (1) it is not appropriate to substitute $\partial G/\partial E$ by $\partial G/\partial V_g$. As a 1D subband is populated there is a tendency for the subband edge to become pinned to the chemical potential, because of the singularity in the 1D density of states $\mathcal{D}(E)$; this has been demonstrated in 1D-2D tunneling,[16] in DC voltage biased conductance measurements,[13] and in capacitance.[17] Although the ob-

served plateaus in conductance and thermopower suggest that such a locking of the Fermi energy may indeed be responsible for the 0.7 structure, the behavior of the conductance with temperature and magnetic field are not easily reconciled with such a model and this may provide a clue to the underlying mechanism.

Pinning of the chemical potential is not the only mechanism which can explain our results. The activated temperature dependence[14] of the 0.7 structure and the enhanced Zeeman splitting[13] suggest a possible spin polarization. If there is a metastable ferromagnetic state which is thermally activated for a significant proportion of the time, electrons of one spin direction will be reflected, and the measured conductance will be a weighted average $\langle G \rangle$ of the transmission properties of the ground state $(2e^2/h)$ and the excited state (e^2/h). This produces a conductance between $2e^2/h$ and e^2/h which explains both the nonquantized value of the 0.7 structure, and the movement of this structure towards e^2/h in a strong parallel magnetic field. The thermopower predicted from Eq. (1) lacks a peak in S corresponding to the transition from $G=0.7(2e^2/h)$ to $G=e^2/h$, but it conflicts with the observations in the region of the nonquantized conductance plateau, as both thermodynamic states would possess conductances independent of energy and zero thermopower; the averaged thermopower will also be zero. An additional mechanism would therefore be required to explain the finite measured thermopower coincident with the 0.7 structure.

Our earlier results[12,13] suggest that the 0.7 structure is accompanied by a spontaneous lifting of the degeneracy between the two spin bands. Increasing this energy difference with a magnetic field reduces the conductance to e^2/h, suggesting that the conductance in excess of e^2/h at $B=0$ may be due to minority spin electrons and the magnetic field reduces the minority spin population. The fact that these minority spins contribute only $0.4(e^2/h)$ to the total conductance might be explained by partial reflection at the entrance to the channel, whereas the majority spin carriers are transmitted normally giving e^2/h. If, contrary to a simple single-electron picture, extra electrons added to the channel enter only the majority spin band, perhaps due to exchange and correlation effects, the predicted behavior would be similar to the case of chemical potential pinning, and a constant conductance but a finite thermopower would result.

In conclusion, the Mott formulation of thermopower due to diffusion of noninteracting electrons can accurately describe ballistic transport through a 1D constriction even in the case of a finite current flow, and a measurement of the differential thermopower allows discrimination between the transmission properties of electrons propagating in opposite directions through the device. We have also investigated the thermopower in the regime of strongly interacting electrons, and although we have not presented a definitive model to explain the finite measured thermopower when the conductance is constant with gate voltage, we have defined the requirements for a successful model.

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[1] N. F. Mott and H. Jones, *The Theory of the Properties of Metals and Alloys*, 1st ed. (Clarendon, Oxford, 1936).
[2] U. Sivan and Y. Imry, Phys. Rev. B **33**, 551 (1986).
[3] P. N. Butcher, J. Phys.: Condens. Matter **2**, 4869 (1990).
[4] C. R. Proetto, Phys. Rev. B **44**, 9096 (1991).
[5] S. F. Godijn, S. Möller, H. Buhmann, L. W. Molenkamp, and S. A. van Langen, Phys. Rev. Lett. **82**, 2927 (1999).
[6] N. J. Appleyard, J. T. Nicholls, M. Y. Simmons, W. R. Tribe, and M. Pepper, Phys. Rev. Lett. **81**, 3491 (1998).
[7] P. J. Price, J. Appl. Phys. **53**, 6863 (1982).
[8] L. W. Molenkamp, H. van Houten, C. W. J. Beenakker, R. Eppenga, and C. T. Foxon, Phys. Rev. Lett. **65**, 1052 (1990).
[9] A. S. Dzurak, C. G. Smith, L. Martín-Moreno, M. Pepper, D. A. Ritchie, G. A. C. Jones, and D. G. Hasko, J. Phys.: Condens. Matter **5**, 8055 (1993).
[10] R. T. Syme, M. J. Kelly, and M. Pepper, J. Phys.: Condens. Matter **1**, L3375 (1989).
[11] N. K. Patel, J. T. Nicholls, L. Martín-Moreno, M. Pepper, J. E. F. Frost, D. A. Ritchie, and G. A. C. Jones, Phys. Rev. B **44**, 13 549 (1991).
[12] K. J. Thomas, J. T. Nicholls, M. Y. Simmons, M. Pepper, D. R. Mace, and D. A. Ritchie, Phys. Rev. Lett. **77**, 135 (1996).
[13] K. J. Thomas, J. T. Nicholls, N. J. Appleyard, M. Y. Simmons, M. Pepper, D. R. Mace, and D. A. Ritchie, Phys. Rev. B **58**, 4846 (1998).
[14] A. Kristensen, P. E. Lindelof, J. B. Jensen, M. Zaffalon, J. Hollingbery, S. W. Pedersen, J. Nygard, H. Bruus, S. M. Reimann, C. B. Sørenson, M. Michel, and A. Forchel, Physica **249-251**, 180 (1998).
[15] A. Yacoby, H. L. Stormer, N. S. Wingreen, L. N. Pfeiffer, K. W. Baldwin, and K. W. West, Phys. Rev. Lett. **77**, 4612 (1996).
[16] L. D. Macks, C. H. W. Barnes, J. T. Nicholls, W. R. Tribe, D. A. Ritchie, P. D. Rose, E. H. Linfield, and M. Pepper, Physica E (Amsterdam) **6**, 518 (2000).
[17] D. Schmerek and W. Hansen, Phys. Rev. B **60**, 4485 (1999).