Interaction-assisted reversal of thermopower with ultracold atoms

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We study thermoelectric currents of neutral, fermionic atoms flowing through a mesoscopic channel connecting a hot and a cold reservoir across the superfluid transition. The thermoelectric response results from a competition between density-driven diffusion from the cold to the hot reservoir and the channel favoring transport of energetic particles from hot to cold. We control the relative strength of both contributions to the thermoelectric response using an external optical potential in both non-interacting and strongly-interacting systems. Without interactions, the magnitude of the particle current can be tuned over a broad range but is restricted to flow from hot to cold in our parameter regime. Strikingly, strong interparticle interactions additionally reverse the direction of the current. We quantitatively model ab initio the non-interacting observations and qualitatively explain the interaction-assisted reversal by the reduction of entropy transport due to pairing correlations. Our work paves the way to studying the coupling of spin and heat in strongly correlated matter using spin-dependent optical techniques with cold atoms.

I. INTRODUCTION

Transport of charge, heat, and spin are often coupled in nature. Their interplay enriches the dynamical response of materials leading to coupled transport phenomena such as thermoelectricity [1] and, along conceptually similar lines, spintronics [2] and spin caloritronics [3]. Thermoelectricity includes two major observations where an applied temperature gradient can induce a charge current (Seebeck effect) or an external voltage can give rise to a heat current (Peltier effect). Besides their widespread practical applications, both are essential to probe fundamental physics. In particular, studies of these effects in strongly correlated materials have allowed researchers to identify relevant charge carriers [4–6] and degrees of freedom [7] which have been proposed to characterize exotic states such as Majorana modes or anyons [8].

Microscopically, thermoelectric currents in conventional materials originate from an electron-hole asymmetry created by an energy-dependent density of states and carrier velocity. In the case of a temperature gradient for instance, this asymmetry can favor the transport of high-energy particles from the hot side over the transport of low-energy particles from the cold side. This imbalance results in a net carrier flow whose magnitude increases with asymmetry.

In solid state systems, several techniques have been explored to engineer the thermoelectric response. First, the energy-dependence of the density of states can be enhanced by reducing the number of free dimensions [9, 10] and using electrostatic gate potentials in low-dimensional structures such as quantum wires [11], point contacts [12] and dots [13]. Furthermore, the thermoelectric response can be strongly modified by electron interactions as observed in quantum dots [13] and two-dimensional electron gases [14]. However, the interpretation of the thermoelectric response in solids is complicated by interactions of the carriers with impurities, defects and phonons.

Due to the absence of these factors, ultracold atoms are well-suited to simulate the relevant physics of real materials. In addition, Feshbach resonances allow one to study the same system across a large range of interaction strengths under comparable conditions. These merits have facilitated experiments on transport phenomena in strongly interacting Fermi gases including spin diffusion [15, 16], sound propagation [17, 18], and heat transport in the form of second sound [19].

As these experiments focused on bulk material properties, they lacked the tunability of mesoscopic systems, however recent work in optically shaping bulk gases made it possible to create mesoscopic cold atom “devices” comparable to their solid state counterparts [20–28]. In particular, thermoelectric phenomena were explored in these mesoscopic structures focusing on either weak [29] or strong [30] interactions. Here, we exploit the ability to compare both interaction regimes in the same structure and, by extending the accessible range of gate potentials, observe a striking reversal of the thermoelectric current directly induced by interactions. This reversal is a novel effect in cold atoms and, to our knowledge, also in strongly correlated solid materials. With the gate, we can finely tune the particle-hole asymmetry at the origin of the thermoelectric current and thereby engineer both its magnitude and direction. We can therefore smoothly turn our system from a heat engine into a heat pump, the latter of which being an important ingredient for an efficient cooling scheme proposed for cold atoms [31].

For weak interactions, the thermoelectric response can be predicted by an ab initio Landauer model thanks to the absence of defects and precise characterization of our system. With strong interactions, we focus on tem-
temperatures around the superfluid transition where a large critical region is predicted \cite{32} and Fermi liquid behavior breaks down \cite{33}. Although the strongly correlated regime is often challenging to understand, we interpret our observation on a fairly fundamental level based on entropy transport. Other works have studied theoretically thermoelectric effects with cold atoms for bosonic \cite{34-38} and fermionic \cite{39-41} systems.

The structure of the paper is as follows. After introducing the setup in Sec. II, we explain its thermoelectric response with an intuitive picture [Sec. III] and discuss the dynamics in the non- and strongly-interacting regimes [Sec. IV]. Based on a phenomenological model presented in Sec. V we extract transport properties and discuss their behaviors in the subsequent Sec. VI and VII. Finally, we conclude in Sec. VIII. Technical details can be found in the Appendices.

II. SETUP

Our transport setup consists of a mesoscopic channel smoothly connected to reservoirs of degenerate fermions (\(^{6}\)Li), as sketched in Fig. 1(a) and described in previous works \cite{21, 42}. The channel is created by a repulsive TEM\(_{01}\)-like laser beam that confines the atoms along the \(z\) direction to its dark nodal plane, reaching a trapping frequency \(\nu_{tx} = 4.5(8)\) kHz at the center. Its Gaussian envelope in the longitudinal direction (\(y\)) ensures the smooth connection of channel and reservoirs. In the \(x\) direction the atoms are restricted by the dipole trap providing a weak confinement with frequency \(\nu_{tx} = 232(1)\) Hz.

The reservoirs, denoted by left (L) and right (R), contain equal atom numbers with \(N = N_L + N_R = 121(2) \times 10^3\) atoms in each of the lowest and third lowest hyperfine states. To prepare a temperature difference we heat one side by parametrically modulating the intensity of an attractive beam inside one reservoir, while blocking the channel. Subsequent equilibration of each reservoir leads to a temperature difference \(\Delta T = T_L - T_R = 147(11)\) nK and an average value \(\bar{T} = (T_L + T_R)/2 = 208(6)\) nK. After reconnecting the reservoirs they exchange particles and heat through transverse modes of the channel.

Their energy is controlled by one of the two gate beams where one is repulsive \((V_g > 0)\) and the other attractive \((V_g < 0)\) [Fig. 1(a)]. By tuning the gate the number of available modes below the chemical potential \(\bar{\mu} = (\mu_L + \mu_R)/2 = 151(16)\) nK reaches up to 40 with 4 occupied states in the tightly confined \(z\) direction making the channel quasi-two-dimensional. Note that non-zero temperature leads to a partial occupation of modes at energies above the chemical potential, which also contribute to transport.

Using a broad Feshbach resonance of \(^{6}\)Li we set the interparticle interactions in the entire system including the channel and the reservoirs either to zero or on resonance (unitarity). Subsequently, we enable transport for a variable time \(t\) and measure the differences in atom numbers \(\Delta N(t)\) and temperature \(\Delta T(t)\) from an absorption image.

III. INTUITIVE LANDAUER PICTURE

In the absence of interactions the response can be understood in a Landauer picture [Fig. 1(b)]. The total thermoelectric current is a result of two competing effects that favor particle currents in opposite directions.
(i) Since the atom numbers are equal in each reservoir the chemical potential difference only depends on the difference in temperatures. Increasing temperature lowers the chemical potential as a result of the particle-hole asymmetry of the reservoir density of states around the Fermi energy. This reservoir asymmetry prefers particle currents from the cold to the hot side. (ii) In order to be transferred from one reservoir to the other, a particle of energy $E$ will transit via one of the transverse modes available in the channel, counted by the transport function $\Phi(E)$. The higher the energy, the more modes are available leading to a faster transfer of energetic particles. This channel asymmetry favors currents from the hot to the cold reservoir, since the hot reservoir has an excess of energetic particles compared to the cold side. The channel asymmetry is controlled by the gate potential inside the channel as intuitively illustrated in Fig. 1(b): For the repulsive case, $V_g > 0$, the channel predominantly allows transport of the most energetic particles present in the hot reservoir. Hence, the net current flows from hot to cold. As the gate potential is made increasingly attractive particles from the cold reservoir start to contribute and reduce the net current until it vanishes. The precise gate potential where the cancellation happens depends on the initial reservoir conditions and on the transport function $\Phi(E)$. At sufficiently attractive gate potentials, $V_g < 0$, almost all particle energies contribute and their direction from cold to hot is purely defined by the chemical potential gradient. We expect the intuitive picture to be also valid for weak interactions as it applies to Landau quasi-particles.

IV. SYSTEM DYNAMICS

Fig. 2(a) presents the experimental evolution of the atom number and temperature difference of a non-interacting Fermi gas subject to an initial temperature gradient for different gate strengths. Due to the combined effect of the asymmetries of the reservoirs and the channel, particles flow from the hot to the cold side and build up a negative difference in atom numbers, $\Delta N = N_L - N_R$. Simultaneously, a heat flow drives the temperatures towards equilibrium thus weakening the thermoelectric current. The particle number difference that builds up induces diffusion that flows against the thermoelectric current and it vanishes at the turning point. Subsequently, diffusion brings the system to equilibrium. This evolution occurs for all gate potentials, except the most attractive one, $V_g = -1.06 \mu K$, where the thermoelectric current vanishes. The observed influence of the gate agrees with the intuitive Landauer picture. For more attractive potentials more transverse modes, counted by the function $\Phi(E)$, open up and permit a faster relaxation. Moreover, the initial response reduces as quantified by the initial particle current $I_N(0)$ in the inset [Appendix B 2]. The current is normalized with the initial bias $\Delta I_0$ to remove variations in the preparation and plotted versus local chemical potential $\mu_{loc}$ at the channel center [Appendix D]. The measurements at strong interactions are presented in Fig. 2(b). As for the non-interacting case the initial temperature bias is converted into a particle number difference and eventually relaxes back to equilibrium. In contrast the initial and relaxation dynamics with interactions are about three times faster and the normalized initial currents are enhanced. This is a consequence of the gate tuning not only the available number of modes but also the density inside the channel. Thus from repulsive to attractive gate potentials the gas becomes denser and is expected to eventually cross the superfluid transition in the channel at $\mu_{loc} = 0.428 \mu K$ (dashed line in inset) [Appendix B 4 and [43]].

Strikingly, near the critical point the net initial current reverses its direction while this effect is absent for an ideal gas in the same parameter regime. In principle the reversal may also be achieved without interactions as predicted from the Landauer picture and is expected to occur for more attractive gate potentials. However, strong attractive gates induce atom losses which prevent the observation of current reversal. Contrary to the present work, the current was flowing only in one direction in previous experiments: from hot to cold for a non-interacting two-dimensional gas in the absence of a gate [29] or from cold to hot in the strongly-interacting, quasi one-dimensional case [30].

As the assumption of a Fermi liquid underlying the Landauer picture breaks down at unitarity, we employ a phenomenological model that captures the transport properties irrespective of the interaction strength. Based on the model we reason in section VI why the reversal occurs with interactions.

V. PHENOMENOLOGICAL MODEL

To extract transport properties from the time evolutions we apply a phenomenological model in linear response. It relates currents of particles $I_N$ and entropy $I_S$ to the differences in chemical potential $\Delta \mu$ and temperature $\Delta T$ [44, 45].

$$
\begin{bmatrix}
I_N \\
I_S
\end{bmatrix}
= G
\begin{bmatrix}
1 & \alpha_c \\
\alpha_c & \alpha_c^2 + L
\end{bmatrix}
\begin{bmatrix}
\Delta \mu \\
\Delta T
\end{bmatrix}
$$

(1)

As a consequence of the Onsager reciprocal relations, the channel properties are described by only three coefficients: the particle conductance $G$, the heat conductance $G_T$, and the Seebeck coefficient or thermopower $\alpha_c$. The thermopower quantifies the coupling between particle and entropy currents and the Lorenz number $L = G_T/TG$ measures the relative strength of particle and heat conductances.

Thermopower is essential to understand how the competing asymmetries of reservoirs and channel appear in the model. At the initial time, the reservoir asymmetry
translates a difference in temperature $\Delta T_0$ to a difference in chemical potential $\Delta \mu = -\alpha_c \Delta T_0$ described by the dilatation coefficient $\alpha_c = - (\partial \mu / \partial T)_N$. Together with Eq. (1) the initial response is

$$I_N(0) = G \alpha \Delta T_0$$

where the effective thermopower $\alpha = \alpha_c - \alpha_r$ determines the outcome of the competition between the asymmetries of the channel and the reservoirs, captured by $\alpha_c$ and $\alpha_r$, respectively.

At strong interactions and lowest temperatures this phenomenological model is expected to fail due to non-linearities caused by superfluid effects in the reservoirs [46]. However at our elevated temperatures it captures the behavior well as observed in [30].

The transfer of particles and entropy between the reservoirs via the associated currents $I_N(t) = - \partial_t \Delta N(t)/2$ and $I_S(t) = - \partial_t \Delta S(t)/2$ modifies the chemical potential and temperature describing each reservoir and in turn the currents. Together with the equation of state in each reservoir the dynamics are captured by a system of coupled differential equations. Their solution governs how the system evolves from the initial conditions $\Delta N_0$ and $\Delta T_0$ towards equilibrium [44] [Appendix B 2],

$$\Delta N(t) = \{e_+(t) + Ae_-(t)\} \Delta N_0 + Be_-(t) \Delta T_0,$$

$$\Delta T(t) = \{e_+(t) - Ae_-(t)\} \Delta T_0 + Ce_-(t) \Delta N_0,$$

with $e_{\pm}(t) = (e^{-t/\tau_+} \pm e^{-t/\tau_-})/2$ and the exponential timescales $\tau_{\pm}$. Both timescales and the coefficients A, B and C depend on the transport properties conductance $G$, thermopower $\alpha_c$ and Lorenz number $L$ and on the thermodynamical properties of the reservoirs through
compressibility, heat capacity and dilatation coefficient \( \alpha_r \). The values of the reservoir properties are taken from the measurement [Appendix B 1 c].

We directly extract the transport coefficients from the transients by fitting \( \Delta N(t) \) and \( \Delta T(t) \) simultaneously, normalized by their statistical uncertainty. To improve accuracy we reduce the number of free parameters by fixing conductance to a separately measured value [Appendix B 3]. The fits are shown as lines in Fig. 2. For the dashed curve conductance was left as a free parameter as it failed to converge otherwise. We discuss the determined transport coefficients in the subsequent chapters.

VI. THERMOPOWER

We first discuss conductance and then present and analyze the results for thermopower. Fig. 3(ab) displays the measured conductances versus local chemical potential \( \mu_{\text{loc}} \) in the non- (dark blue) and strongly-interacting (orange) regime. Conductance is increased by interactions, as previously observed [42] and is a factor of 13(1) larger at the most attractive gate strength [Fig. 3(b)]. In the non-interacting regime the solid line represents an ab initio prediction that reproduces the measured well. The model is based on Landauer’s theory that follows the idea that both reservoirs inject particles according to their occupations and the channel transmits them in one of the available modes [47] [Fig. 1(b) and Appendix E]. The light blue point indicates the fitted conductance corresponding to the dashed curve in Fig. 2(a) and is consistently higher than model and separate measurement.

The behavior of the thermopower is summarized in Fig. 3(c). In the non-interacting regime the thermopower \( \alpha_c \) of the channel (blue points) reduces with increasing local chemical potential \( \mu_{\text{loc}} \) (horizontal blue line). As expected the reduction originates from a suppression of the channel asymmetry with more attractive gate strengths as illustrated in Landauer’s picture. Moreover, the dilatation coefficient is a property of the reservoirs and thus unaffected by the gate that locally acts on the channel. The Landauer prediction is indicated as a solid black curve and reproduces well the extracted thermopower in the non-interacting case, except for the most attractive gate strength (\( \mu_{\text{loc}} = 1.1 \mu\text{K} \)). There theoretically the reversal of the thermoelectric current is anticipated (\( \alpha_c < \alpha_r \)), while experimentally it is absent (\( \alpha_c \gtrsim \alpha_r \)) which might originate from the detailed shape of the confining potential not captured in the theory, such as anharmonicities. Overall, the tunability of the thermopower in the non-interacting regime demonstrates almost full control of the thermoelectric response. In contrast, at unitarity the thermopower \( \alpha_c \) becomes smaller than the dilatation coefficient \( \alpha_r \) for sufficiently attractive gate potentials and additionally \( \alpha_r \) and \( \alpha_c \) are reduced compared to the non-interacting case.

The interaction-assisted reduction and reversal can be qualitatively understood from an interpretation based on entropy. On the one hand, the role of the reservoirs is captured by the dilatation coefficient that can be expressed as the entropy content to add a particle isothermally, \( \alpha_c = (\partial S/\partial N)_T \). On the other hand, rewriting Eq. (1) as

\[
I_S = \alpha_c I_N + G_T \frac{\Delta T}{T} \tag{5}
\]

allows for reinterpreting the thermopower of the channel as the average entropy that is reversibly transported by one particle, while the second term captures the irreversible entropy exchange between the reservoirs. At unitarity, pairing correlations reduce the entropy in the spin sector and account for the decrease of \( \alpha_c \) and \( \alpha_r \) compared to the non-interacting case, as visible in Fig. 3(c) and in [43]. Because the attractive gate increases the density and therefore the interaction effects at the center, we expect further reduction of the thermopower \( \alpha_r \).

This argument suggests a way to estimate where the current reverses in the presence of interactions. Since the interparticle collision rate is enhanced at unitarity, we assume the gas to be locally in equilibrium at the center [Appendix C]. This allows us to think of the thermopower \( \alpha_r \) as the dilatation coefficient inside the channel at temperature \( T \) and chemical potential \( \mu_{\text{loc}} \) leading to a reversal expected at \( \mu_{\text{loc}} \sim -0.1 \mu\text{K} \). The value is relatively close to the observed location despite the simplicity of the estimation and the neglect of the transverse mode structure.

VII. LORENZ NUMBER

The Lorenz number \( L = G_T/TG \) compares the ability of systems to conduct heat and particles and indicates whether Fermi liquid behavior is present. In this case transport is described by Landau quasiparticles that carry both charge and energy leading to a constant Lorenz number, \( L_{\text{WF}} = \pi^2/3 \cdot k_B^2 \). This is known as the Wiedemann-Franz law.

Fig. 4 displays the fitted Lorenz number versus local chemical potential in the non- (blue dots) and strongly-interacting (orange dots) case. Without interactions the extracted values are consistently higher than \( L_{\text{WF}} \) and also than the Landauer theory (solid line). It considers our mesoscopic geometry at finite temperature and approaches \( L_{\text{WF}} \) with increasing degeneracy. Deviations from the Wiedemann-Franz law are found in other systems with either increased [48] or decreased [49–51] numbers. Here, we partly attribute the inconsistency to small systematic shifts in the measured temperature difference that mostly affect the Lorenz number [Appendix B 3].

Despite the challenging absolute estimation the Lorenz number is reduced by one order of magnitude when increasing the interactions to unitarity. As the fitted heat conductance is relatively insensitive to interaction strength, this decrease can be mostly attributed to the
enhancement of the conductance seen in Fig. 3(ab). Also the smaller uncertainties on the Lorenz number stem from the enhanced conductances and give us confidence in thinking that the Wiedemann-Franz law is violated here, as experimentally observed in a one-dimensional geometry [30] and theoretically supported in [40, 41].

VIII. DISCUSSION

Mesoscopic transport properties are influenced by the geometry of the structure [52] thus it is instructive to compare our findings at unitarity with a one-dimensional channel measured in similar conditions [30]. There, the narrow geometry blocked irreversible heat currents leading to a non-equilibrium steady state with a finite temperature difference. Relaxation was restored by widening the channel thanks to an enhanced heat exchange. This agrees with our results for a wide quasi-two-dimensional channel where we observe relaxation to equilibrium and a heat conductance that is similar to the non-interacting case. In both works the Lorenz number is reduced with interactions by an order of magnitude. However the reasons are different: In [30] it stems from a reduced heat conductance while here it is almost purely an effect of enhanced particle conductance. Thermopower is reduced by interactions in our wide quasi two-dimensional geometry as qualitatively explained by pairing correlations that restrict average entropy transfer per particle. In contrast, in the quasi one-dimensional channel it shows the surprising behavior of following the non-interacting prediction, which so far eluded any explanation [40].

In summary, we control magnitude and direction of thermoelectric currents through a mesoscopic structure in the presence of weak and strong interparticle interactions. In our parameter regime at weak interactions particles are flowing consistently in one direction while at unitarity we observe a striking interaction-assisted reversal. We explain the reversal by a competition of reservoir and channel parameters which are affected by pairing correlations that are expected in the large critical region around the superfluid transition [32]. Indeed, the reversal occurs before the normal-to-superfluid transition and its precise location depends on the geometry of the system and the reservoirs conditions via thermopower and dilatation coefficient. Unitary Fermi gases are not the only system where thermopower can be affected by interactions: In a two-dimensional Bose gas the Seebeck coefficient changes its sign close to the superfluid transition [53] and in two-dimensional electron gases interactions are predicted to reduce the thermopower and potentially reverse its sign [54].

The option to induce thermoelectric currents in either direction is appealing when considering our dynamics as...
an open thermodynamic cycle. In our system, an atomic flow from hot to cold acts against the chemical potential bias and converts heat into work as a thermoelectric engine. Conversely, the system acts as a thermoelectric cooler when the flow transfers heat from cold to hot. The initial direction of the current therefore determines which mode of operation takes place before the other: a thermoelectric engine in the channel-dominated regime or a cooler in the reservoir-dominated regime. In both modes the conversion efficiency of the cycles is characterized by the figure of merit $ZT = \frac{\alpha^2 T}{L} [45]$. Contrary to a non-interacting system interactions strongly reduce the Lorentz number while thermopower remains at a similar order of magnitude. Overall, we estimate that interactions improve the figure of merit by a factor of 7(3) showing the relevance of strongly correlated quantum materials for thermoelectric applications.

Our system readily allows to probe the thermoelectric response of more complicated structures. Drawing from the technique to imprint local effective Zeeman shifts [55] it opens new perspectives on coupling spin and heat transport [3]. By adding strong correlations intriguing thermoelectric effects could be observed [56, 57].

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Appendix A: Experimental details

A cloud of fermionic lithium atoms is prepared in a balanced mixture of the lowest and third lowest hyperfine state in a hybrid trap. In the transverse directions $(x, z)$ the trap is formed by a far-detuned laser beam at a wavelength of 1064 nm and longitudinally $(y)$ by a quadratic magnetic field. We evaporatively cool the elongated cloud by reducing the transverse confinement frequency $\nu_{tx} = 232(1)$ Hz and $\nu_{tz} = 212(1)$ Hz in the transverse direction and longitudinally $\nu_{ty}$ of 26.6(2) Hz at non-interacting and 32.2(2) Hz at unitarity.

The transport channel is imprinted by a repulsive laser beam at 532 nm that is shaped by a 0–π phase plate into a TEM$_{00}$ mode and focused onto the center of the cloud. This compresses the cloud in vertical $z$ direction resulting in a harmonic confinement of $\nu_{cz} = 4.5(8)$ kHz. Its $1/\epsilon^2$ waists are in longitudinal direction $w_{cy} = 30.2\mu m$ and transversally $w_{cz} = 9.5\mu m$.

We prepare the initial conditions of the two reservoirs in two steps. First, the cloud is centered on the transport channel using a magnetic gradient to make the atom numbers on both sides equal, and then we split the reservoirs with an elliptical repulsive laser beam. Second, a temperature difference is created by heating one side with a red-detuned beam at 767 nm focused into one reservoir and modulating its intensity parametrically. The experimentally optimized modulation frequency is 547 Hz, which is on the order of the transverse confinement frequencies $\nu_{tx}$ and $\nu_{tz}$. The beam position is controlled with a piezo-mirror and can be directed into either reservoir. After letting the reservoirs thermalize for 10 ms we reach the initial conditions reported in Table I.

The attractive gate potential is created with the same beam used for heating and requires a careful alignment onto the channel center. Its waists are $w_{gx} = 34.3(2)\mu m$ and $w_{gy} = 33.5(2)\mu m$. The repulsive one is formed by an elliptical beam at 532 nm that is focused onto our channel with waist $w_{gx} = 53.66(3) \mu m$ in transverse direction and along transport $w_{gy} = 8.58(3)\mu m$. Besides acting as a gate the same beam enables and blocks transport at large powers in a controlled way. After the time $t$ the reservoirs are again separated and an absorption picture is taken after a short time of flight of 1 ms. This reduces the central densities and allows to image with intensities well below saturation. From the absorption image we extract the thermodynamical properties as described in B1c.

### TABLE I. Reservoir conditions for measuring thermoelectricity in the non- and strongly-interacting regimes. The values are averaged over the realizations at different gate potentials and their uncertainties indicate one standard deviation.

| Quantity | Non-interacting | Unitary |
|----------|-----------------|---------|
| $N$ (10$^5$) | 121(2) | 105(4) |
| $T_L$ (nK) | 281(8) | 226(5) |
| $T_R$ (nK) | 314(7) | 129(3) |
| $\mu_L$ ($k_B$ nK) | -43(27) | 65(13) |
| $\mu_R$ ($k_B$ nK) | 344(12) | 238(8) |
| $T$ (nK) | 208(6) | 174(2) |
| $\bar{\mu}$ ($k_B$ nK) | 151(16) | 152(9) |
| $\Delta T$ (nK) | 147(11) | 90(7) |
| $\Delta \mu$ ($k_B$ nK) | -387(29) | -172(13) |
Appendix B: Data analysis

1. Thermodynamics

Across the appendix, the equation of state of a non-interacting and unitary Fermi gas are used in the homogeneous and trapped cases. This section summarizes the relevant thermodynamical relations.

a. Homogeneous unitary Fermi gas

The universality hypothesis states that the homogeneous unitary Fermi gas is described by the interatomic distance and the thermal wavelength only [58]. This restricts the equation of state to the form [59]

\[
n \lambda_T^3 = f_n(q)
\]

with the particle density \(n\), the thermal wavelength \(\lambda_T = \sqrt{2\pi^2/mk_BT}\) and the dimensionless function \(f_n\) depending on the reduced chemical potential \(q = \beta \mu\). The scaling function \(f_n\) was measured around the superfluid transition [43] and theoretically extended in the degenerate and thermal limits with a phonon model and a third order virial expansion, respectively [59][60]. Using the equation of state the normalized temperature is given as

\[
T/T_F = \left(\frac{4}{3\sqrt{\pi}f_n(q)}\right)^{2/3}
\]

with the Fermi energy \(E_F = k_B T_F = \hbar^2/(2m)(6\pi^2n)^{2/3}\) [59]. This links the normalized temperature \(T/T_F\) and the chemical potential \(\mu/E_F = T/T_F \cdot q\) to each other. Below the critical temperature \(T_c = 0.167 \cdot T_F\) the gas becomes superfluid as observed in [43].

the compressibility \(\kappa\), the specific heat \(C_N\) and the internal energy \(U\) are relevant. Following reference [61] they are given in the non-interacting case by

\[
N = \left(\frac{k_B T}{\hbar \nu}\right)^3 F_2(q_0),
\]

\[
\frac{\kappa}{N} = \frac{1}{2} \frac{F_1(q_0)}{F_2(q_0)},
\]

\[
\frac{\alpha_r}{k_B} = 3 \frac{F_2(q_0)}{F_1(q_0)} - q_0,
\]

\[
\frac{C_N}{Nk_B} = 12 \frac{F_3(q_0)}{F_2(q_0)} - 9 \frac{F_2(q_0)}{F_1(q_0)},
\]

\[
U = \frac{3}{6^{1/3}} \frac{F_3(q_0)}{F_2(q_0)^{5/3}},
\]

\[
\frac{T}{T_F} = \frac{4}{3\sqrt{\pi}f_n(q_0)}^{1/3},
\]

with \(F_j\) symbolizing the complete Fermi-Dirac integral of order \(j\) [62, Eq. (25.12.14)]. At unitarity the relations are stated in [63] and read as,

\[
N = \frac{4}{\sqrt{\pi}} \left(\frac{k_B T}{\hbar \nu}\right)^3 N_2(q_0),
\]

\[
\frac{\kappa}{N} = \frac{1}{2k_B T} N_0(q_0),
\]

\[
\frac{\alpha_r}{k_B} = \frac{6}{N_0(q_0)} - q_0,
\]

\[
\frac{C_N}{Nk_B} = \frac{8}{3} \frac{N_4(q_0)}{N_2(q_0)} - \frac{18}{N_0(q_0)},
\]

\[
U = \frac{1}{\sqrt{\pi}} N_2(q_0)^{3/2},
\]

\[
\frac{T}{T_F} = \left(\frac{24}{\sqrt{\pi}} N_2(q_0)\right)^{-1/3},
\]

with the integral \(N_j(q) = \int_0^\infty r^j f_n(q-r^2) \, dr\) involving the scaling function \(f_n\) defined in section B1a.

To describe the properties of a single reservoir we need to divide the extensive quantities for the full trap by two. Here, this concerns the atom number \(N\), the compressibility \(\kappa\), the specific heat \(C_N\) and the internal energy \(U\).

c. Extracting thermodynamical properties

The basis to extract thermodynamical properties of each reservoir \(r\) from the absorption images is the virial theorem [64]. It holds for non-interacting and unitary [65] Fermi gases and relates the internal energy per particle \(U_r/N_r = 3\hbar^2 \nu_r/2m\) to the second moment \(\langle g^2 \rangle = \int g^2 n_r(y) \, dy \, / \, N_r\) in transport direction. The one-dimensional density \(n_r(y)\) can be directly deduced from the absorption images by summing over the transverse direction. Together with the total atom number \(N_r\) we obtain the Fermi energy \(E_{F,r}\) and subsequently with Eq. (B7) or (B13) the reduced chemical potential \(\beta \mu_0\).
at the trap center. Then, the temperature $T_r$ and all other thermodynamical properties of the reservoirs can be derived from the formulae in appendix B 1 b.

2. Parameter extraction

The following section explains first the linear phenomenological model [44] and second its application to extract the transport parameters from the measured evolutions [Fig. 2].

a. Phenomenological model

The evolutions in the differences in atom number and temperature are phenomenologically modeled by the response matrices of the channel and the reservoirs. The channel reacts to the applied biases with currents of particles and entropy [Eq. (1)] described by

$$\begin{align*}
I_N &= G \left( \frac{1}{\alpha_c} - \frac{L}{\alpha_r^2} \right) \cdot \frac{\Delta \mu}{\Delta T}, \\
I_S &= \frac{\kappa}{\alpha_r} \cdot \frac{\Delta \mu}{\Delta T}.
\end{align*}$$

(B15)

with the transport coefficients conductance $G$, thermopower $\alpha_c$ and Lorenz number $L$. The reservoir thermodynamics are described in linear response by

$$
\begin{align*}
\Delta N &= \kappa \left( \frac{1}{\alpha_r} - \frac{L}{\alpha_r^2} \right) \cdot \frac{\Delta \mu}{\Delta T}.
\end{align*}
$$

(B16)

The involved thermodynamical quantities of each reservoir are the compressibility $\kappa$, the dilatation coefficient $\alpha_r$ and the reservoir analogue of the Lorenz number $l = C_N/T\kappa$ that depends on the heat capacity $C_N$. Within linear response they are constant throughout the evolution and are evaluated at the equilibrium chemical potential $\bar{\mu}$ and temperature $\bar{T}$ [Appendix B 1 b]. The two sets of equations are related via $I_N(t) = -\partial_t \Delta N(t)/2$ and $I_S(t) = -\partial_t \Delta S(t)/2$ that describe that changes in the reservoirs originate from currents leaving and entering them. Altogether, they form a system of two coupled linear differential equations with the solution given by

$$
\begin{align*}
\Delta N(t) &= \{ c_+(t) + Ae_-(t) \} \Delta N_0 + Be_-(t) \Delta T_0, \\
\Delta T(t) &= \{ c_+(t) - Ae_-(t) \} \Delta T_0 + Ce_-(t) \Delta N_0,
\end{align*}
$$

(B17)

with $c_\pm(t) = (e^{-t/\tau} \pm e^{-t/\tau-})/2$ and the timescales $\tau_\pm = \tau_0/\lambda_\pm$ and $\tau_0 = \kappa/2G$. The eigenvalues of the evolution matrix describing the system of differential equations are

$$
\lambda_\pm = \frac{1}{2} \left( 1 + \frac{L + \alpha^2}{l} \right) \pm \sqrt{\frac{1}{4} \left( 1 + \frac{L + \alpha^2}{l} \right)^2 - \frac{L}{l}}.
$$

(B19)

The coefficients $A$, $B$ and $C$ depend on the transport and reservoir properties

$$
\begin{align*}
A &= \frac{1 - (L + \alpha^2)/l}{\lambda_+ - \lambda_-}, \\
B &= \frac{2\kappa\alpha}{\lambda_+ - \lambda_-}, \\
C &= \frac{2\alpha}{\kappa l(\lambda_+ - \lambda_-)}
\end{align*}
$$

(B20)

(B21)

(B22)

with the effective thermopower $\alpha = \alpha_c - \alpha_r$. Note that these equations correct typographical errors present in [44].

b. Data fitting

We extract the transport parameters by fitting simultaneously the measured evolutions to the phenomenological model using the standard least squares method. It minimizes the following sum of the squared residuals:

$$
\chi^2 = \sum_i \left[ \frac{(\Delta \tilde{N}(t_i) - \Delta \tilde{N}_i)}{\sigma_N} \right]^2 + \left[ \frac{(\Delta \tilde{T}(t_i) - \Delta \tilde{T}_i)}{\sigma_T} \right]^2
$$

(B23)

The simultaneous fitting requires normalizing the residuals with the respective uncertainties $\sigma_N$ and $\sigma_T$ which we determine as an average over the standard deviations at the individual times. The measured differences $\Delta \tilde{N}_i$ and $\Delta \tilde{T}_i$ at the time $t_i$ consist of small offsets due to imperfections in the calibrations. We account for these by shifting the model evolutions $\Delta N(t)$ and $\Delta T(t)$ by constants and obtain $\Delta \tilde{N}(t) = \Delta N(t) + \Delta N_{\text{off}}$ and $\Delta \tilde{T}(t) = \Delta T(t) + \Delta T_{\text{off}}$, respectively.

In the model the reservoir properties $\kappa$, $\alpha_r$ and $l$ are set to the values directly extracted from the absorption images [Appendix B 1 c]. Moreover, the conductance $G$ is fixed by a separate measurement [Appendix B 3] which reduces the number of free parameters and improves the fit.

We determine $\Delta N_0$ and $\Delta T_0$ from the corresponding measured initial values. The offset $\Delta N_{\text{off}}$ follows either from an average over points where the evolution is relaxed or is taken to be the same as the initial value. In contrast, the offset $\Delta T_{\text{off}}$ cannot be easily determined for evolutions that do not completely relax within our measurement. Thus we fit it consistently for all dynamics and checked with evolutions where temperature completely relaxes that the results do not depend on whether the offset $\Delta T_{\text{off}}$ is fixed or left free. In summary the parameters $\alpha_c$, $L$ and $\Delta T_{\text{off}}$ are free in the model. Exceptionally, for the non-interacting curve at $V_g = 0.37$ mK conductance is additionally left free as otherwise the fit failed to converge [Fig. 2].

To analyze the parameter estimation we exemplify the sum of the squared residuals $\chi^2$ versus thermopower $\alpha$ and Lorenz number $L$ for the gate potential $V_g = 0.13$ mK.
For both interaction strengths there exists a single, isolated global minimum that is found by the Levenberg-Marquardt algorithm (gray point). Its uncertainty only includes the standard deviations $\sigma_N$ and $\sigma_T$ and assumes the fixed parameters to be precisely known. However, their uncertainties will lead to a change in the optimal values for $\alpha$ and $L$. Fig. 5(ab) shows a plot of these optimal values when one fixed parameter is varied within its uncertainty. These variations are mostly symmetrical and thus leave the values unaffected but increase the overall error bar in thermopower and Lorenz number.

It is known that estimating parameters of exponential models is challenging as small changes in the data can strongly influence the parameters [66]. Thus in addition to the least squares method we used a technique based on rank order that is robust against outliers and shown to improve the estimation for exponential models [67, 68]. Applying this method led to comparable results and hence we conclude that for our model and data the least squares method performs well. As the estimation with the rank order fit is numerically demanding due to many local closely spaced minima [67] we report the values with the least squares method which gives a single well-defined minimum.

Based on the fitted results the initial response $I_N(0)/\Delta T_0 = G\alpha$ is shown in the insets of Fig. 2, the thermopower of the channel $\alpha_c = \alpha + \alpha_r$ in Fig. 3 and the Lorenz number $L$ in Fig. 4.

3. Conductance measurement

To measure conductances we prepare the reservoirs at different particle numbers and equal temperatures and study its exponential decay towards equilibrium. We fit the characteristic time $\tau_0$ and together with the compressibility $\kappa$ of each reservoir [details in appendix B 1] conductance follows from the relation $\tau_0 = \kappa/2G$ analogous to a discharging capacitor [24]. The conductances are measured at the same average chemical potential $\mu$ and temperature $T$ as the thermoelectric responses to allow comparison between them. The conductances are shown in Fig. 3(ab).

The conductances are extracted under the assumption that both reservoirs are at equal temperatures. However, we suspect that in our measured temperature differences an offset is present, as visible in Fig. 2 when the system is relaxed. During preparation we heat both reservoirs up to seemingly the same temperature $T$ while the measurement offset leads to a physical temperature difference. In turn it induces a thermoelectric current with value

$$I_N = G \frac{\Delta N}{\kappa} \left( 1 + \kappa \alpha \frac{\Delta T}{\Delta N} \right), \tag{B24}$$

which follows from Eq. (B15) and (B16) and the thermopower $\alpha = \alpha_c - \alpha_r$. The second summand in the brackets quantifies the relative deviations that are around 10% estimated for the non-interacting case at a gate potential $V_g$ of 0.13 mK. We used the compressibility $\kappa = 22 s/h$, the thermopower $\alpha = 0.7k_B$, the particle number difference $\Delta N = 48 \cdot 10^3$ and the residual temperature bias $\Delta T \sim -15 nK$.

A systematically increased conductance mostly reduces the fitted Lorenz number and only slightly decreases the thermopower, as visible in Fig. 5(cd). The Lorenz number is reduced to around $4.2k_B^2$ for a gate potential of $V_g = 0.13 mK$ and similarly for the other non-interacting conditions. Thus the systematic shifts in the conductance can partly explain the deviations from the Wiedemann-Franz value $L_{WF} = \pi^2/3 \cdot k_B^2$ in the non-interacting measurements.

4. Superfluid transition at unitarity

Superfluidity inside the reservoirs or the channel can strongly influence the transport properties and thus it is important to characterize the gas in these locations in the strongly-interacting regime.

The transition is captured by the thermodynamics of the gas and characterized by the degeneracy $T/T_F$ as presented in Fig. 6 versus the local chemical potential. At the channel center (black points) the gas is in a hydrodynamic regime thanks to the strong interparticle interactions [Appendix C] and thus it is represented by the local equilibrium described by the average temperature $T = 172 nK$ and the local chemical potential $\mu_{loc}$ [Appendix D]. Thus with increasing chemical potential the degeneracy increases as stated by the thermodynamics of a homogeneous unitary Fermi gas [Appendix B 1 a] and eventually crosses the critical value $T_c/T_F = 0.167(13)$ at $\mu_{loc} = 0.428 \, \mu K$ (gray dotted lines) [43].

The hot (red points) and cold (blue points) reservoirs are not affected by the gate potential and hence their degeneracy remains constant. They are defined by a half-harmonic potential where the densest point at the center specifies the transition at a critical value $T_c/T_F = 0.223$ (horizontal violet dotted line). Note that the value depends on the geometry via the Fermi temperature $T_F$. In summary the reservoirs are both above the critical temperature while the channel center transitions from a non-condensed to a superfluid state depending on the chemical potential. Thus in our system the gas may become superfluid inside the channel.

Appendix C: Boltzmann approach at unitarity

Inside the channel interatomic collisions may alter the transport regime from ballistic to hydrodynamic depending on how frequently they scatter while crossing the constriction. To assess the regime at unitarity we follow the Boltzmann approach outlined in [69] and compare the resulting mean free path with the length of the channel. First we focus on the center where the gas is homogeneous and then extend the argument for the full channel.
FIG. 5. Transport parameter estimation. (ab) Sum of squared residuals $\chi^2$ versus effective thermopower $\alpha$ and Lorenz number $L$ in the (a) non- and (b) strongly-interacting regimes at $V_g = 0.13\mu K$. Contour lines highlight the behavior at logarithmically spaced level values indicated by vertical black lines in the colorbar. The offset $\Delta T_{off}$ in temperature difference is fixed to the fitted value. The optimal fit parameters and its uncertainties are indicated by the point and the error bar, respectively. (cd) Variations of the optimal thermopower and Lorenz number when one fixed fit parameter is changed within its uncertainty (colored lines). Left panel: Modifications in conductance (orange) and compressibility (blue). Right panel: Modifications in initial values $\Delta N_0$ (solid green), $\Delta T_0$ (solid violet) and the offsets $\Delta N_{off}$ (dashed green) and $\Delta T_{off}$ (dashed violet). Changes in the thermodynamic analogue of the Lorenz number $l$ are too small to be visible.

1. Scattering time

Within the Boltzmann theory the scattering rate $\gamma$ follows from the collision integral and reads [69]

$$\gamma N = \frac{m}{\pi^2 \hbar^3} \int dE_1 dE_2 dE_3 dE_4 \delta(E_1 + E_2 - E_3 - E_4) \\ \times (1 - f_1)(1 - f_2)f_3 f_4 \cdot H(E, E_m), \quad (C1)$$

where $N$ denotes the number of particles in each hyperfine state. The integral describes two particles with energies $E_3$ and $E_4$ that exchange energy during the collision and leave with energies $E_1$ and $E_2$. As the collision is elastic the total particle energy $E = E_1 + E_2 = E_3 + E_4$ is conserved, as visible from the delta function. The product of the Fermi-Dirac distributions $f_i = 1/(1 + e^{\epsilon_i})$ with $\epsilon_i = (E_i - \mu)/k_B T$ ensures that the initial states are occupied and the final ones unoccupied, as dictated by Pauli’s principle. The factor $H$ is the integrated density of states weighted by the collision cross section $\sigma(k)$. In our situation the gas is locally homogeneous at the center which leads to

$$H(E, E_m) = \frac{2\pi m V}{(2\pi\hbar)^3} \int_{P_-}^{P_+} dP \sigma(k), \quad (C2)$$

with the volume $V$ and the integral over total momentum $P$ whose bounds are $P_k = \sqrt{2m(E - E_m)} \pm \sqrt{2mE_m}$ with the minimal energy $E_m = \min(E_1, E_2, E_3, E_4)$. At unitarity the collision cross section is $\sigma(k) = 4\pi/k^2$ expressed with the relative wavevector $k$ given by $4(hk)^2 = P_x^2 + P_y^2 - P_z^2$. 

$$
\begin{align*}
\Delta\alpha (k_B) & = 0.05 \\
\Delta L (k_B) & = 0.1 \\
\Delta N_0 (k_B) & = -0.07 \\
\Delta T_0 (k_B) & = 0.07 \\
\end{align*}
$$
For the prefactor we choose the classical collision rate in \( \gamma \) from the variables smallest and that the total energy is conserved we transform the variables \( E \) to the non-negative \( n \) and \( x \) and normalize with the Fermi energy \( E_F = \hbar^2/(2m)(6\pi^2n)^{2/3} \) with the particle density \( n = N/V \).

\[
\begin{align*}
E_1/E_F &= x_1, \\
E_2/E_F &= x_1 + x_2, \\
E_3/E_F &= x_1 + x_3, \\
E_4/E_F &= x_1 + x_2 + x_3.
\end{align*}
\]

The resulting scattering rate \( \gamma = \gamma_0 I(T/T_F) \) consists of a prefactor \( \gamma_0 \) and a dimensionless integral \( I(T/T_F) \). For the prefactor we choose the classical collision rate in a homogeneous gas at temperature \( T_F \)

\[
\gamma_0 = n \cdot \sqrt{2} \tilde{v} \cdot \sigma(kF) = \frac{8\sqrt{2}}{3\pi^{3/2}} \frac{E_F}{\hbar},
\]

with the average relative velocity \( \sqrt{2}\) between two particles expressed through their individual mean speed \( \bar{v} = \sqrt{8E_F/\pi m} \). The dimensionless integral reads

\[
I \left( \frac{T}{T_F} \right) = \frac{9\sqrt{\pi}}{4} \int_0^\infty dx_1dx_2dx_3 f(x_1 + x_2)f(x_1 + x_3) \\
\times [1 - f(x_1)][1 - f(x_1 + x_2 + x_3)] \\
\times F(2x_1 + x_2 + x_3, x_1),
\]

with \( f(x) = 1/(1 + e^x) \) and \( \xi = (T_F/T)(x - \mu/E_F) \) where the normalized temperature \( T/T_F \) and chemical potential \( \mu/E_F \) are related in appendix B 1a. The function \( F(x, x_m) \) depends on the normalized total \( x = E/E_F \) and minimal energy \( x_m = x_1 \) as

\[
F(x, x_m) = \frac{1}{\sqrt{x}} \log \left( \frac{2 + x/x_m + 2\sqrt{2x/x_m}}{2 + x/x_m - 2\sqrt{2x/x_m}} \right). \tag{C9}
\]

This form is unsuited for numerical evaluation as it contains a divergence at \( x = 0 \) which we lift with the substitution \( w = x_1, y = (x_2 + x_3)/2x_1 \) and \( z = (x_3 - x_2)/2x_1 \) leading to the final result

\[
I \left( \frac{T}{T_F} \right) = \frac{9}{2} \sqrt{\pi} \int_0^\infty dy \log \left( \frac{x + y/2 + \sqrt{1 + y}}{x + y/2 - \sqrt{1 + y}} \right) \\
\times \int_0^\infty dw [1 - f(w)][1 - f(w(1 + 2y))] \sqrt{\frac{w^3}{1 + y}} \\
\times \int_{-y}^y dz f(w(1 + y - z))f(w(1 + y + z)). \tag{C10}
\]

Instead of scattering rates we use its inverse, the average time between interparticle collisions \( \tau = \tau_0/I(T/T_F) \).

2. Mean free path

Based on the average scattering time calculated in the previous section we estimate the mean free path \( l_{mfp} = \tau v_F \) a particle with velocity \( v_F = \sqrt{2E_F/m} \) travels between collisions. For the estimation we need the local degeneracy \( T/T_F \) and the Fermi energy \( E_F \) at the center. They both follow from the temperature \( T \) and chemical potential \( \mu \) via the local chemical potential \( \mu_{loc} \), as detailed in appendix D, and the thermodynamics of a homogeneous unitary Fermi gas in section B 1a.

3. Transport regime

To discuss the transport regime we first focus on the scattering time and mean free path at the channel center [Fig. 7]. Versus local chemical potential the classical scattering time \( \tau_0 \) (dashed line) monotonically decreases which indicates more frequent collisions. Although the cross section \( \sigma(kF) \sim 1/E_F \) reduces as particles collide at higher relative velocities the simultaneous increase in density \( n \sim E_F^{1/2} \) and mean particle speed \( \bar{v} \sim E_F^{1/2} \) dominates [compare with Eq. (C7)]. In the quantum case (solid line) the time \( \tau \) behaves the same at small chemical potentials. In contrast at larger values the gas is more degenerate and Pauli’s principle limits collisions as it requires the final states to be empty. Eventually this effect dominates and leads to a longer time between collisions. The local minimum in scattering time appears less pronounced in the mean free path as the Fermi velocity monotonically increases [Fig. 7(b)].
The predictions based on Boltzmann’s approach are valid in the normal Fermi liquid phase located above the critical and superfluid regimes. In the critical region pairing correlations are relevant and modify the scattering rate. In a harmonic trap they were found to almost compensate Pauli blocking giving rates that follow the classical prediction [70]. Hence, we expect in the homogeneous case that in the critical region Pauli blocking is compensated giving rates that follow the superfluid regime.

Appendix D: Effective potential

In two-terminal setups particles move along one direction and are confined in the others. As a result of the confinement the transverse motion is quantized into different modes labelled by quantum numbers. As long as the confinement varies adiabatically through the channel particles remain in the same mode which is called adiabatic approximation. Then, the confinement energy acts as an additional potential as it is invested in the transverse direction and is missing for the longitudinal motion. In the following we present the effective potential of our channel that is useful to calculate the local chemical potential at the center and to deduce transport parameters [appendix E].

Harmonically approximating our transverse confinement \((x, z)\) at each longitudinal position \((y)\) leads to the potential

\[
V(x, y, z) = \frac{1}{2} m \omega_x^2(y) x^2 + \frac{1}{2} m \omega_z^2(y) z^2 + V_g(y). \tag{D1}
\]

The harmonic frequencies \(\omega_x(y)\) and \(\omega_z(y)\) include a nearly constant contribution from the dipole trap due to the long Rayleigh length of 20 mm and a spatially varying one from the gate potential and the one creating the channel, respectively. They are given by

\[
\begin{align*}
\omega_x^2(y) &= \omega_{x0}^2 + \omega_{gg}^2 e^{-2y^2/w_{gy}^2} \tag{D2} \\
\omega_z^2(y) &= \omega_{z0}^2 + \omega_{gg}^2 e^{-2y^2/w_{gy}^2} \tag{D3}
\end{align*}
\]

with the dipole trap frequencies \(\omega_{tx}/\omega_{tz} = 2\pi \cdot \nu_{tx}/\nu_{tz}\), the frequencies at the center created by the gate \(\omega_{gx}^2 = -4V_g/m\omega_{gx}^2\) and channel beam \(\omega_{cz}\). Their waists in longitudinal direction are \(w_{gy}\) of 8.6 \(\mu\)m for the repulsive and 33.5 \(\mu\)m for the attractive gate beam and \(w_{cy} = 30.2\) \(\mu\)m for the channel beam. Besides modifying the trapping frequency the gate beam creates an additional potential

\[
V_g(y) = V_g e^{-2y^2/w_{gy}^2}. \tag{D4}
\]

As at each longitudinal position \(y\) the transverse potential is quadratic and shifted by the energy \(V_g(y)\) the eigenenergies are

\[
E_n(y) = \hbar \omega_x(y) (n_x + 1/2) + \hbar \omega_z(y) (n_z + 1/2) + V_g(y) \tag{D5}
\]

with the quantum number \(n = (n_x, n_z)\). Fig. 8 displays the eigenenergies in transport direction for different gate strengths \(V_g\).

1. Local chemical potential

To characterize the gas at the channel center we use the local density approximation with the effective potential \(V_{\text{eff}}(V_g) = E_0(0)\) in the ground state. This results in a local chemical potential \(\mu_{\text{loc}} = \bar{\mu} - V_{\text{eff}}(V_g)\) with \(V_{\text{eff}}(V_g) = \hbar (\omega_x(0) + \omega_z(0))/2 + V_g\). Note that \(\omega_z(0)\) implicitly depends on \(V_g\) in a square root fashion. The local chemical potential is used to display transport coefficients [Fig. 3 and 4], to locate the superfluid transition [section B-4] and to discuss its scattering properties [appendix C].

![FIG. 7. Boltzmann approach for a homogeneous unitary Fermi gas at the channel center. (a) Interparticle scattering time in classical (dashed curve) and quantum (solid line) regimes and (b) mean free path versus local chemical potential. The horizontal line indicates the channel length \(2w_{cy}\) in terms of the waist \(w_{cy}\) of the creating laser beam. The superfluid transition (vertical dashed line) is estimated in appendix B4. The predictions are valid in the normal regime above the critical region. In the critical region we expect the real values to be lower (see text) and in the superfluid regime the theory fails.](image-url)
2. Transport function $\Phi(E)$

From the reservoirs to the center the channel narrows and tends to increase the mode energies, visible in Fig. 8 in the absence of the gate ($V_g = 0$). Here, a particle at energy $E$ (horizontal line) can cross the channel in any mode indicated in blue. An additional repulsive potential ($V_g > 0$) pushes them up and further peaks them at the center within the size of the beam. In the attractive case ($V_g < 0$) the energies are pulled down and some modes might be energetically allowed at the center while away from it they are above (red lines). Only the modes whose energies are below the particle energy throughout the channel are relevant for transport. Their number is counted with the transport function $\Phi(E)$ directly from the effective potential and is used in appendix E to calculate transport parameters.

Appendix E: Landauer-Büttiker theory

In this appendix we detail how we model the transport parameters in the non-interacting situation. In the Landauer framework particles come from the reservoirs following its Fermi-Dirac distribution and cross the channel with a transmission probability. In linear response the transport coefficients follow and are indicated in Fig. 3 and 4.

$$G = \frac{1}{\hbar} \int_{-\infty}^{+\infty} \Phi(E) \left(-\frac{\partial f(\epsilon)}{\partial E}\right) dE,$$

$$G\alpha_c = \frac{1}{\hbar T} \int_{-\infty}^{+\infty} \Phi(E)(E - \mu) \left(-\frac{\partial f(\epsilon)}{\partial E}\right) dE,$$

$$G(L + \alpha^2) = \frac{1}{\hbar^2 T} \int_{-\infty}^{+\infty} \Phi(E)(E - \mu)^2 \left(-\frac{\partial f(\epsilon)}{\partial E}\right) dE,$$

with the Fermi-Dirac distribution $f(\epsilon) = 1/(1 + e^{\epsilon/(k_B T)})$ and the normalized particle energy $\epsilon = (E - \mu)/k_B T$.

In the classical regime the transmission $\Phi(E)$ through the channel reduces to the number of transverse modes below energy $E$. We count their number directly from the effective potential discussed in appendix D. Then by numerically evaluating the Landauer integrals the transport coefficients follow and are indicated in Fig. 3 and 4.

1. Benchmarking

To benchmark the method we compare it with the measured conductance for a non-interacting gas shown in Fig. 3(ab) and find good agreement. Note that counting available modes at the channel center leads to a wrong prediction that increases roughly quadratically with local chemical potential, in contrast to the observed linear behavior.

2. Validity of linear response

In Landauer theory particle and entropy currents are expressed with the difference $\Delta f(E) = f_L(E) - f_R(E)$ between the Fermi-Dirac distributions. To arrive at the linear response form [Eq. (E1), (E2) and (E3)] the difference is expanded around the average chemical potential $\overline{\mu}$ and temperature $\overline{T}$ and we obtain

$$\Delta f(E) \approx -\frac{\partial f(\epsilon)}{\partial E} \left(\Delta \mu + \epsilon k_B \Delta T\right)$$

with the derivative $-\partial f(\epsilon)/\partial E = 1/(2k_B T)[1 + \cosh(\epsilon)]$. The distributions and the exact and linearized differences are displayed in Fig. 9 for the conditions of the non-interacting measurement. Visibly linearization only introduces minor deviations justifying the approximation.

[1] J. He and T. M. Tritt, Science 357 (2017).
[2] I. Žutić, J. Fabian, and S. Das Sarma, Rev. Mod. Phys. 76, 323 (2004).
[3] G. E. W. Bauer, E. Saitoh, and B. J. van Wees, Nature Materials 11, 391 EP (2012).
[4] J. M. Toumazou, Journal of Physics: Condensed Matter 30, 183001 (2018).
[5] O. Cyr-Choinière, S. Badoux, G. Grisonnanche, B. Michon, S. A. A. Afshar, S. Fortier, D. LeBoeuf, D. Graf, J. Day, D. A. Bonn, W. N. Hardy, R. Liang, N. Doiron-Leyraud, and L. Taillefer, Phys. Rev. X 7, 031042 (2017).
[6] F. Laliberté, J. Chang, N. Doiron-Leyraud, E. Hassinger, R. Daou, M. Rondeau, B. J. Ramshaw, R. Liang, D. A. Bonn, W. N. Hardy, S. Pyon, T. Takayama, H. Takagi, I. Sheikin, L. Malone, C. Proust, K. Behnia, and L. Taillefer, Nature Communications 2, 432 (2011).
[7] Y. Wang, N. S. Rogado, R. J. Cava, and N. P. Ong, Nature 423, 425 (2003).
[8] Y. Kleeorin, H. Thierschmann, H. Buhmann, A. Georges, L. W. Molenkamp, and Y. Meir, Nature Communications 10, 5801 (2019).
[9] M. S. Dresselhaus, G. Dresselhaus, X. Sun, Z. Zhang, S. B. Cronin, and T. Koga, Physics of the Solid State 41, 679 (1999).
[10] M. Dresselhaus, G. Chen, M. Tang, R. Yang, H. Lee, D. Wang, Z. Ren, J.-P. Fleurial, and P. Gogna, Advanced Materials 19, 1043 (2007).
[11] Y. Dubi and M. Di Ventra, Rev. Mod. Phys. 83, 131 (2011).
[12] H. van Houten, L. W. Molenkamp, C. W. J. Beenakker, and C. T. Foxon, Semiconductor Science and Technology 7, B215 (1992).
FIG. 8. Eigenenergies $E_n$ with transverse mode $n$ versus longitudinal position $y$ (a) with repulsive, (b) without and (c) with attractive gate potential $V_g$. Modes available throughout the channel at energy $E$ (horizontal line) are indicated in blue and contribute to transport. The ones in red are accessible close to the center but blocked away from it and the remaining modes are plotted in gray.

FIG. 9. Validity of linear response. (a) The Fermi-Dirac distributions in both reservoirs as prepared for the non-interacting measurement and (b) their exact (solid) and linearized (dashed) difference. The used numbers are summarized in Table I.

[13] A. A. M. Staring, L. W. Molenkamp, B. W. Alphenaar, H. van Houten, O. J. A. Buyk, M. A. A. Mabesoone, C. W. J. Beenakker, and C. T. Foxon, EPL (Europhysics Letters) 22, 57 (1993).
[14] A. Makashi, S. Li, B. Wen, S. V. Kravchenko, A. A. Shashkin, V. T. Dolgopolov, and M. P. Sarachik, Phys. Rev. Lett. 109, 096405 (2012).
[15] A. Sommer, M. Ku, G. Roati, and M. W. Zwierlein, Nature 472, 201 EP (2011).
[16] M. Koschorreck, D. Pertot, E. Vogt, and M. Köhl, Nature Physics 9, 405 EP (2013).
[17] P. B. Patel, Z. Yan, B. Mukherjee, R. J. Fletcher, J. Struck, and M. W. Zwierlein, arXiv e-prints , arXiv:1909.02555 (2019), arXiv:1909.02555 [cond-mat.quant-gas].
[18] M. Bohlen, L. Sobirey, N. Luick, H. Biss, T. Enss, T. Lonpe, and H. Moritz, arXiv e-prints , arXiv:2003.02713 (2020), arXiv:2003.02713 [cond-mat.quant-gas].
[19] L. A. Sidorenkov, M. K. Tey, R. Grimm, Y.-H. Hou, L. Pitaevskii, and S. Stringari, Nature 498, 78 EP (2013).
[20] M. Albiez, R. Gati, J. Fölling, S. Hunsmann, M. Cristiani, and M. K. Oberthaler, Phys. Rev. Lett. 95, 010402 (2005).
[21] J.-P. Brantut, J. Meineke, D. Stadler, S. Krinner, and T. Esslinger, Science 337, 1069 (2012).
[22] S. Eckel, J. G. Lee, F. Jendrzejewski, N. Murray, C. W. Clark, C. J. Lobb, W. D. Phillips, M. Edwards, and G. K. Campbell, Nature 506, 200 (2014).
[23] F. Jendrzejewski, S. Eckel, N. Murray, C. Lanier, M. Edwards, C. J. Lobb, and G. K. Campbell, Phys. Rev. Lett. 113, 045305 (2014).
[24] S. Krinner, D. Stadler, D. Husmann, J.-P. Brantut, and T. Esslinger, Nature 517, 64 (2015).
[25] G. Valtolina, A. Burchianti, A. Amico, E. Neri, K. Xhani, J. A. Seman, A. Trombettoni, A. Smerzi, M. Zaccanti, M. Inguscio, and G. Roati, Science 350, 1505 (2015).
[26] C.-C. Chien, S. Peotta, and M. Di Ventra, Nat Phys 11, 998 (2015).
[27] S. Krinner, T. Esslinger, and J.-P. Brantut, Journal of Physics: Condensed Matter 29, 343003 (2017).
[28] L. Amico, M. Boshier, G. Birkl, A. Minguzzi, C. Miniatura, L. C. Kwek, D. Aghamalyan, V. Ahufinger, N. Andrei, A. S. Arnold, M. Baker, T. A. Bell, T. Bland, J. P. Brantut, D. Cassettari, F. Chevy, R. Citro, S. De Palo, R. Dumke, M. Edwards, R. Follman, J. Fortagh, S. A. Gardiner, B. M. Garraway, G. Gauthier, A. Günter, T. Haug, C. Hufnagel, M. Keil, W. von Klitzing, P. Ireland, M. Lebrat, W. Li, L. Longchambon, J. Mompert, O. Morsch, P. Nakdesi, T. W. Neely, M. Olshanii, E. Orignac, S. Pandey, A. Pérez-Obiol, H. Perrin, L. Pirola, J. Polo, A. L. Pritchard, N. P. Proukakis, C. Ry-
