Superfluid current through a dissipative quantum point contact

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We measure superfluid transport of strongly-interacting fermionic lithium atoms through a quantum point contact with local, spin-dependent particle loss. We observe that the characteristic non-Ohmic superfluid transport enabled by high-order multiple Andreev reflections survives even at dissipation strength greater than the superfluid gap. We develop a model with mean-field reservoirs connected via tunneling to a dissipative site. Our calculations in the Keldysh formalism reproduce the observed non-equilibrium particle current, yet do not fully explain the observed loss rate or spin current.

The interplay between coherent Hamiltonian dynamics and the incoherent, dissipative dynamics emerging from coupling to the environment is essential to the physics of open quantum systems [1–3]. This leads to rich phenomena including the quantum Zeno effect [4–8], emergent dynamics [9–14], and dissipative phase transitions [15–20]. Furthermore, an important question is how many-body coherence competes with dissipation through dephasing or particle loss. Directed transport between two reservoirs offers an advantageous geometry for studying this competition, since dissipation can be applied locally without perturbing the many-body states robustly provided by the reservoirs [21]. So far, studies on dissipation in solid-state systems have focused on dephasing [22, 23]. More recently, quantum gases have been established as a versatile platform to study interacting many-body physics and to engineer novel forms of dissipation [2, 4], though previous experiments on dissipation in transport settings have focused on weakly interacting systems [24–26].

Engineered dissipation in strongly-correlated fermionic systems, while less explored [19, 27], is particularly interesting as such systems exhibit high-order processes relying on many-body coherence, which one might expect to be strongly suppressed by dissipation. A prime example is the excess current between two superconductors [28] or superfluids [29] through a quantum point contact (QPC). Due to the superfluid gap ∆, quasiparticles cannot be transported alone from one reservoir to another when the chemical potential difference ∆μ between the reservoirs is smaller than 2∆ [illustrated in Fig. 1(d)]. Instead, transport occurs via simultaneous co-tunnelling of quasiparticles and nMAR = ∆/∆μ Cooper pairs [30, 31], a process known as multiple Andreev reflection (MAR) [32–35]. The robustness of such high-order processes to dissipation is an interesting open question, especially for pair-breaking particle loss acting on only one spin state, since the very existence of MAR relies on many-body coherence between the spins.

In this work, we address this question by experimentally and theoretically studying the influence of spin-dependent particle loss on superfluid transport. We use a strongly-correlated Fermi gas—a superfluid with many-body pairing—in a transport setup with two reservoirs connected by a QPC and apply controllable local particle loss at the QPC with dissipation strengths up to several times ∆—the energy scale responsible for the superfluid current. We find that, surprisingly, the superfluid behavior survives. This result is reproduced by a minimal model which includes both superconductivity and dissipation written in the Keldysh formalism.

**Experiment.**—We prepare a degenerate Fermi gas of 6Li in a harmonic trap in a balanced mixture of the first- and third-lowest hyperfine ground states, labeled |↓⟩ and |↑⟩. The atomic cloud has typical total atom numbers \( N = N_↑ + N_↓ = 195(14) \times 10^3 \), temperatures \( T = 100(2) \) nK, and Fermi temperatures \( T_F = 391(10) \) nK. Using a pair of repulsive, TEM₉₀₁-like beams intersecting at the center of the cloud, we optically define two half-harmonic resonance frequencies \( \nu_z \approx 11(2) \) kHz and \( \nu_z \approx 9.9(2) \) kHz, realizing a QPC illustrated in Fig. 1(a). We apply a magnetic field of 689.7 G to address the spins’ Feshbach resonance, giving rise to a fermionic superfluid in the densest parts of the cloud at the contacts to the 1D channel. An attractive gaussian beam propagating along \( z \) acts as a “gate” potential \( V_g \) which increases the local chemical potential at the contacts, which determines both the superfluid gap \( ∆ \) and the number \( n_m \) of occupied transverse transport modes in the 1D region. Both \( ∆ \) and \( n_m \) can be computed from the known potential energy landscape and equation of state [36] and are ap-
approximately $\Delta/k_B \approx 1.4 \mu K$ and $n_m \approx 3$ in this work.

We engineer spin-dependent particle loss with a tightly focused beam at the center of the 1D channel [Fig. 1(a)] that optically pumps $|\downarrow\rangle$ to an auxiliary ground state $|5\rangle$ [Fig. 1(c)] which interacts weakly with the two spin states and is lost due to photon recoil [36]. This leads to a controllable particle dissipation rate of $|\downarrow\rangle$ atoms given by the peak photon scattering rate $\Gamma_\downarrow$ with no observable heating. While this loss beam is far off-resonant for $|\uparrow\rangle$, strong interactions lead to loss of $|\uparrow\rangle$ [36]. Even at the strongest dissipation explored here, the system lifetime is over a second, much longer than the timescale of a detectable transport of $10^3$ atoms via MAR—about $10^3 h/\Delta \sim 5$ ms.

We induce particle transport from the left to the right reservoir by preparing an atom number imbalance $\Delta N = N_L - N_R$ which generates a chemical potential bias $\Delta \mu = \Delta \mu(\Delta N, N, T)$ given by the system’s equation of state [36]. $\Delta \mu$ drives a current $I_N = -\Delta N/2$ from left to right that causes $\Delta N$ to decay over time $t$. The dynamics of $\Delta N(t)/N(t)$ for various $\Gamma_\uparrow$ are plotted in Fig. 2(a). From each trace, we numerically extract the current-bias relation $I_N(\Delta \mu)$, shown in Fig. 2(b) in units of the superfluid gap $\Delta$ [36].

For weakly-interacting spins, the decay of $\Delta N(t)$ at arbitrary $\Gamma_\uparrow$ is exponential since a QPC coupling two Fermi liquids has a linear (Ohmic) current-bias relation $I_N = G\Delta \mu$. The conductance $G$ is quantized in units of the inverse Planck constant $2/h$ (2 comes from spin degeneracy) but can be renormalized to a smaller value by dissipation [25, 46]. In contrast, we observe that the decay of $\Delta N$ at $\Gamma_\uparrow = 0$ deviates strongly from an exponential and the corresponding current-bias relation $I_N(\Delta \mu)$ is highly non-linear, consistent with previous observations in the strongly-interacting regime [29]. In fact, this non-linearity is a signature of superfluidity. Specifically, for ballistic QPCs such as ours, the current is approximately $I_N \approx G\Delta \mu + I_N^{\text{qc}}$ where the normal current $G\Delta \mu \approx 2n_m \Delta \mu/h$ is carried by quasiparticles and the excess current $I_N^{\text{qc}} \approx (16/3) n_m \Delta h$ is carried by the Cooper pairs participating in MAR and is directly given by the superfluid gap $\Delta$—the natural energy scale in this system [34]. The dominance of the excess current over the normal current can be seen directly in $\Delta N(t)/N(t)$ with $\Gamma_\uparrow = 0$: It decays almost linearly in time, indicating that the current is nearly independent of $\Delta \mu$.

We now turn to the question of the influence of dissipation on the superfluid transport. At low $\Gamma_\uparrow$, superfluidity is still evident from the linear initial decay of $\Delta N$ which gives way to exponential behavior at low bias where the MARs responsible for transport become higher order. The effect of dissipation is clearer in the extracted current-bias relations shown in Fig. 2(b). As dissipation strength increases, the initial current (largest $\Delta \mu$) reduces as does the concavity of the curve. The current-bias relation eventually approaches a straight line (Ohmic) at high dissipation, yet the slope is still significantly larger than the normal conductance $2n_m/h$. This is similar to the effect of increasing temperature, where the excess current decreases due to a reduced $\Delta$ but superfluid fluctuations, even in the normal phase, lead to anomalously large conductance [29, 61, 64–66].

**Theoretical model and fit.**—We have developed a mean-field model based on previous work [29] by adding an atom loss process within the Lindblad master equation framework [36]. The two reservoirs are treated as BCS superfluids, and transport through the channel is modeled by single-particle tunneling with amplitude $\tau$ from either reservoir onto a dissipative site between them [Fig. 1(b)]. The dissipation is modeled with a Lindblad operator $\hat{L}_\sigma = \sqrt{\gamma_\sigma} d_\sigma$ proportional to the site’s fermionic annihilation operator $d_\sigma$ with a rate $\gamma_\sigma$ for each spin $\sigma = \downarrow, \uparrow$, i.e. as a pure particle loss process that is uncorrelated between the two spins. We include $\gamma_\downarrow$ as a phenomenological dissipation to account for the loss of $|\downarrow\rangle$ atoms due to strong pairing correlations.

To compute nonequilibrium observables such as $I_N$, we use the Keldysh formalism extended to dissipative systems [19, 38–42, 67]. The time integral of the theoretical current $I_N^\text{K}(\Delta \mu, \tau, \gamma_\downarrow, \gamma_\uparrow)$ along with the equation of state $\Delta \mu(\Delta N, N, T)$ yields the model’s prediction for the time evolution of the particle imbalance $\Delta N^\text{K}(t, \tau, \gamma_\downarrow, \gamma_\uparrow)$. Here, $N(t)$ is obtained from an exponential fit to the data [36]. For simplicity, we model a single transport mode and obtain the net current by multiplying by the number of modes $\Delta \dot{N}^\text{K} = -2n_m I_N^\text{K}$. Although the formalism can treat non-zero temperatures,
we simplify the calculation of $I_N^K$ by using zero temperature since $k_B T / \Delta < 0.08$.

Due to the spin-dependent loss, a spin imbalance builds up, leading to a magnetization imbalance $\Delta M = \Delta N_1 - \Delta N_0 \neq 0$ which can drive additional particle current. Nevertheless, $\Delta M$ remains small ($\Delta M / N < 0.07$) for all our data and its effect on $I_N^K$ is negligible [36]. Moreover, the model predicts a spin current one order of magnitude below the observed value, and therefore does not fully describe the spin-dependent transport. In this work, we focus on spin-averaged particle transport.

We perform a least-squares fit of $\Delta N^K(t, \tau, \gamma, \gamma_0)$ to the measured $\Delta N(t)$ to extract the model parameters $\tau$, $\gamma$, and $\gamma_0$. In our case of low bias and ballistic channel, $I_N^K$ is sensitive only to the sum $\gamma = \gamma_0 + \gamma$. Therefore, to avoid overfitting, we fix the ratio $\gamma_0 / \gamma = r$ to the average measured ratio of the atom loss rates of the two spin states $r = N_0 / N_1 = 0.72(4)$ [36]. As the gap $\Delta$ is the natural unit of current and bias, the estimation of $\Delta$ in the experimental system is crucial for quantitative comparison to the model. However, it is not clear how the spatially varying gap in the potential energy landscape of the experiment (crossover from the 3D reservoirs to the 1D channel) enters the model. Motivated by this and an experimental uncertainty of about 6% in the potential energy of the gate beam, we fit a multiplicative correction factor $\eta_g$ on the gate potential $V_g$, which strongly affects both our estimate of $\Delta$ in the most degenerate point in the system and $n_m$ [36]. Since the system is identical in each data set except for the dissipation strength, we first fit the $\Gamma_1 = 0$ data with $\gamma_0 = \gamma = 0$ to find $\tau$ and $\eta_g$, which determine the concavity of $I_N^K(\Delta \mu)$ and the timescale of $\Delta N^K(t)$, respectively. We then fit the single parameter $\gamma_0$ with fixed $r$, $\tau$, and $\eta_g$ for subsequent sets. The systematic error in $\eta_g$ is the dominant source of uncertainty in our theoretical calculations.

The fits for each data set are plotted as solid lines in Fig. 2(a,b). With this fitting procedure, we find $\eta_g = 1.00(6)$, $\Delta / k_B = 1.4(1) \mu K$, and $n_m = 2.7(4)$—all errors include experimental systematic uncertainties. The fitted $\tau$ can be represented in physical units by the energy linewidth $\Gamma_d$ of the dissipative site as $\Gamma_d \propto \tau^2$ [36] and gives $\Gamma_d = 5.9(4) \Delta$. Independent measurements with different values of $\nu_s$ and $V_g$ produce similar results for $\eta_g$ and $\Gamma_d$ within 20%. The large value of $\Gamma_d$ reflects the near-perfect transmission of the ballistic channel as the limit $\Gamma_d \to \infty$ is equivalent to perfect transparency $\alpha \to 1$ in a QPC model with direct tunneling between the reservoirs [29, 36, 37]. In principle, the energy $\epsilon_d$ of the dissipative site is an additional free parameter. Due to the large linewidth, however, the model is insensitive to changes of $\epsilon_d$ within the physically meaningful range $|\epsilon_d| < \Delta$. There are thus no resonant effects as in a weakly-coupled quantum dot, and we fix $\epsilon_d = 0$. For the same reason, we do not consider any on-site interaction.

The fitted $\gamma$ and $\gamma_0$ versus $\Gamma_1$ plotted in the inset of Fig. 2(a), show that the effective dissipation strength is approximately proportional to the photon scattering rate as expected. The fitted slope $\gamma_0 = 3.61(13) \Gamma_1$ is of order 1 since the dissipative beam’s waist $w_g = 1.31(2) \mu m$ is comparable to the Fermi wavelength in the channel $\lambda_F \approx 2 \mu m$ and thus shows that the single-site model captures the physics of the real situation of a continuum of states in the channel and strong interaction.

Robustness of superfluid transport to dissipation.—To see how the superfluid character—the excess current—changes with dissipation strength quantitatively, we replot the measured and calculated current in Fig. 3(a) versus $\Gamma_1$ at a few bias values including the initial bias $\Delta \mu / \Delta \approx 0.05$ ($n_{\text{MAR}} \approx 20$) down to $\Delta \mu / \Delta \approx 0.005$ ($n_{\text{MAR}} \approx 200$). The fitted model is shown in solid curves using the fitted linear relationship between $\gamma_0$ and $\Gamma_1$ in the inset of Fig. 2(a). The upper bound of the normal current from the quantized conductance of quasiparticles is $2n_m \Delta \mu / h$ (dashed lines in corresponding colors). The observed current well exceeds this bound for all biases, indicating the persistence of the large excess cur-
FIG. 3. Robustness of high-order MAR to dissipation. (a) Measured current vs. $\Gamma_1$ at different biases. The theoretical calculations are shown as solid curves with uncertainties in lighter colors. Dashed lines represent upper bounds of normal-state currents $2n_m\Delta \mu/h$. (b) Same data at $\Delta \mu/\Delta \approx 0.05$ with bounds on the conserved current $I_{\text{cons}}^\text{upper} = [I_N - \tilde{N}/2, I_N + \tilde{N}/2]$ indicated by the shaded area (lighter color represents the uncertainty). The horizontal axis is the fitted $\gamma$ in units of the gap to illustrate the effective strength of the dissipation. The conserved current is above the normal current (bounds represented by the horizontal bar), showing superfluid character even for the strongest dissipation.

rent. Moreover, the excess current decays smoothly as dissipation increases and gives no indication of a dissipative phase transition but rather shows a dissipative superfluid-to-normal crossover (cf. Ref.[6]).

Since the current $I_N = -(\tilde{N}_L - \tilde{N}_R)/2$ could in principle arise purely from an asymmetric particle loss, we verify that the conserved current $I_{\text{cons}}^\text{upper}$—the atoms transported through the QPC without being lost—exceeds the upper bound for the normal current for dissipation strengths well above the superfluid gap. This is shown in Fig. 3(b) where we plot the data for the largest bias together with bounds for $I_{\text{cons}}^\text{lower}$. These bounds are obtained by writing $\tilde{N}_L/R = \pm I_{\text{cons}}^\text{upper} - I_{L/R}^\text{loss}$, where $I_{L/R}^\text{loss}$ is the dissipation-induced current into the vacuum [42], and thus $I_N = I_{\text{cons}}^\text{lower} + (I_{L/R}^\text{loss} - I_{R}^\text{loss})/2$. Assuming the worst-case scenario—maximally asymmetric loss—leads to $I_{\text{cons}}^\text{lower} \in [I_N - \tilde{N}/2, I_N + \tilde{N}/2]$. These results show that the system preserves its superfluid nature, and the high order processes responsible for the large excess current have partially survived in the measured regime.

Atom loss rate. Finally, we compare the theoretical model to the experiment in terms of the total atom loss rate $\dot{N} = \dot{N}_L + \dot{N}_R$. In the Keldysh formalism, computing steady-state observables involves an integration over energy. Because our model assumes reservoirs with linear dispersion, a high-energy cutoff $\Lambda$ is needed to obtain realistic results for some observables, effectively incorporating a minimal form of energy dependence into the model [36]. The loss rate in the model is determined by $\gamma_{\text{1,1}}$, the tunneling amplitude $\tau$, and the cutoff $\Lambda$. The fit to the $\dot{N}(t)$ data determines $\tau$ and $\gamma_{\text{1,1}}$, as explained above while $\Lambda$ only affects the total atom loss rate $\dot{N}$ and has no influence on the calculated current as long as $\Lambda > \Delta$. Intuitively, the Fermi energy $E_F \approx 2.3\Delta$ [48] is a reasonable energy limit, though $\Lambda$ should be larger than the linewidth of the dissipative site $\Gamma_\text{d} \approx 5.8\Delta$, which is another relevant energy scale. We find that a single $\Lambda$ cannot reproduce the observed atom loss rate at different dissipation [shown in Fig. 4(a) for $\Delta \mu/\Delta \approx 0.05$]: There is a saturation of atom loss rate as dissipation increases, while the calculated loss rate increases almost linearly with dissipation in the measured range. In particular, $\Lambda \approx 5\Delta$ is needed to reproduce the data at large dissipation and $\Lambda \gtrsim 20\Delta$ for the low dissipation data. This discrepancy is shown in Fig. 4(b) which plots the $\Lambda$ needed to reproduce the atom loss rate at each measured $\Gamma_1$ in Fig. 4(a). The calculations do predict a saturation of the atom loss rate at an order of magnitude higher dissipation strength followed by a slow decrease of $\dot{N}$ versus $\Gamma_1$—a signature of the quantum Zeno effect [4–6]. We do not observe such an effect here, yet it cannot be excluded at stronger dissipation. However, the lack of energy resolution in $\dot{N}$ [68] and the finite width of the dissipative beam are expected to wash out its signature.

We attribute the discrepancies between theory and experiment to the highly simplified treatment of the reservoirs as BCS superfluids with linear dispersion and the lack of energy dependence in the tunneling and dissipation strengths. In reality, we work in the unitary regime where pairing correlations and fluctuations are important [48]. The energy dependence in the reservoirs and channel neglected in this model can be important as it is responsible for the strong thermoelectric effect present in this system [44, 69]. This leaves open questions for future investigations of a more refined model.

Conclusion.—We have shown that the superfluid current of a strongly-correlated Fermi gas through a QPC survives spin-dependent particle loss with dissipation.
strengths larger than the superfluid gap, demonstrating a remarkable robustness of the high-order, coherent processes responsible for the current. We observed no critical behaviour, which stands in stark contrast to other types of superfluid-suppressing perturbations such as moving defects [70–72], magnetic fields in solid-state superconductors [73], and photon absorption by superconducting nanowire single-photon detectors [74]. While our most significant observations are captured by our minimal model, deviations from the theory point to the importance of pairing correlations and energy dependence of the dissipation. These effects can be probed directly in future studies on the thermoelectricity [44], spin conductance [64, 75], and correlated loss [76–79] in a similar system.

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SUPPLEMENTAL MATERIAL

I. SUMMARY OF THE THEORETICAL CALCULATION

A. Lossy quantum dot coupled to reservoirs

We model the channel between the reservoirs by a lossy quantum dot coupled to leads, described as an open quantum system for which the dynamics of the density operator $\hat{\rho}$ obeys the Lindblad master equation [1]

$$\frac{d\hat{\rho}}{dt} = -i[\hat{H}, \hat{\rho}] + \sum_{\sigma=1,\downarrow} \gamma_{\sigma} \left( \hat{d}_{\sigma} \hat{\rho} \hat{d}_{\sigma}^\dagger - \frac{1}{2} \left( \hat{d}_{\sigma}^\dagger \hat{d}_{\sigma} \hat{\rho} + \hat{\rho} \hat{d}_{\sigma}^\dagger \hat{d}_{\sigma} \right) \right).$$

Here, $\gamma_{\sigma}$ is the dissipation rate of spin state $\sigma$ and $\hat{d}_{\sigma}$ ($\hat{d}_{\sigma}^\dagger$) the fermionic annihilation (creation) operator at the quantum dot. Notice also that we adopt natural units $\hbar = k_B = 1$. The loss within the channel is taken into account by a local particle loss at the quantum dot. The Hamiltonian is $\hat{H} = \hat{H}_L + \hat{H}_R + \hat{H}_d$, where $\hat{H}_L$ and $\hat{H}_R$ describe the superfluid reservoirs, $\hat{H}_d$ the quantum dot, and $\hat{H}_t$ the tunneling between the reservoirs and the dot.

We treat effects of the superfluid reservoirs within the mean-field Hamiltonian (see Ref. [29])

$$\hat{H}_t = \sum_k \hat{\psi}_{ik}^\dagger \left[ (\epsilon_k - \mu_i) \sigma_z + \Delta \sigma_x \right] \hat{\psi}_{ik}$$

where $i = L, R$ for left and right and $\epsilon_k$ represents the single-particle energy at momentum $k$. As in the case of Ref. [29], we consider a linear single-particle dispersion relation $\epsilon_k = hv_F(k - k_F)$, corresponding to a constant density of states in the normal state as opposed to the quadratic density of states for harmonically-trapped gases. The chemical potential in each reservoir is $\mu_i$, $\Delta$ is the superfluid energy gap, and $\sigma_{x,z}$ are Pauli spin matrices. Here, $\hat{\Psi}_{ik}$ denotes the Nambu spinor $\hat{\Psi}_{ik} = (\hat{\psi}_{ik↑}, \hat{\psi}_{ik↓})^T$, and $\hat{\psi}_{ikσ}$ ($\hat{\psi}_{ikσ}$) is the fermionic creation (annihilation) operator for reservoir $i$. A current through the quantum dot is induced by a chemical potential difference $\Delta\mu = \mu_L - \mu_R$ between the two reservoirs. We choose the chemical potentials symmetrically as $\mu_L = \Delta\mu/2$ and $\mu_R = -\Delta\mu/2$. While the reservoirs are interacting, we consider noninteracting fermions at the quantum dot, with the Hamiltonian $\hat{H}_d = \epsilon_d \sum_{\sigma=1,\downarrow} \hat{d}_{\sigma}^\dagger \hat{d}_{\sigma}$. Here, $\epsilon_d$ is the quantum dot energy level. Single-particle tunneling occurs between the position $r = 0$ in each reservoir and the quantum dot:

$$\hat{H}_t = -\tau \sum_{\sigma=1,\downarrow} \left[ \hat{\psi}^\dagger_{Lσ}(0) \hat{d}_{σ} + \hat{d}_{σ}^\dagger \hat{\psi}_{Rσ}(0) + \text{h.c.} \right].$$

The tunneling amplitude $\tau$ is energy-independent and is a fitted parameter as are the dissipation rates $\gamma_{\sigma}$. Note that $\tau$ has units of (energy)$^{-1}$ (length)$^{3/2}$ for 3D reservoirs. Its corresponding energy scale, the linewidth or inverse lifetime of the quantum dot, is given by $\Gamma_d = \pi \rho_0 \tau^2$ where $\rho_0$ is the density of state of the reservoirs at the Fermi level in the normal state. The quantum dot energy level is set to $\epsilon_d = 0$ in the middle transparency of the reservoir chemical potentials, where the high transparency is realized [37].

Transport is connected to the change in particle numbers of the reservoirs $\hat{N}_i = \hat{N}_{i↑} + \hat{N}_{i↓}$, where $\hat{N}_{iσ} = \int d\mathbf{r} \hat{\psi}^\dagger_{iσ}(\mathbf{r}) \hat{\psi}_{iσ}(\mathbf{r})$. For an open quantum system described by the quantum master equation (1), the time derivative of the particle number is obtained as (see Ref. [38])

$$\frac{d}{dt} \langle \hat{N}_i \rangle = \frac{d}{dt} \text{Tr} \left( \hat{N}_i \hat{\rho}(t) \right) = i\tau \sum_{\sigma=1,\downarrow} \left( \langle \hat{\psi}^\dagger_{iσ}(0) \hat{d}_{σ} \rangle - \langle \hat{d}_{σ}^\dagger \hat{\psi}_{iσ}(0) \rangle \right).$$

The apparent current observed in the experiment is given by $I_N = -\frac{1}{2\pi} \text{Tr} (\hat{N}_L - \hat{N}_R)$ and the loss rate is $\dot{N} = \frac{1}{2\pi} (\hat{N}_L + \hat{N}_R) = -\sum_{\sigma=1,\downarrow} \gamma_{\sigma} \langle \hat{d}_{σ}^\dagger \hat{d}_{σ} \rangle$.

B. Keldysh formalism

We apply the Keldysh formalism [39, 40] to compute the nonequilibrium expectation values, which can be expressed in terms of Keldysh Green’s functions. By using the path integral formulation, the Keldysh action is written in the basis of fermionic coherent states parametrized by the Grassmann variables $\psi$. Integration over a closed time contour is performed by introducing $\psi = (\psi^+, \psi^-)$ for the forward and backward time branches. For the convenience of calculation, we perform the following change of the field variables [39]:

$$\begin{pmatrix} \psi^1 \\ \psi^2 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & 1 \\ 1 & -1 \end{pmatrix} \begin{pmatrix} \psi^+ \\ \psi^- \end{pmatrix},$$

which allows to express the expectation values in terms of advanced, retarded, and Keldysh components of Green’s function. As we are interested in equal-time correlations, it is also convenient to use the frequency basis. The relevant expectation values only depend on $\psi$ and $\bar{\psi}$ at $\mathbf{r} = 0$, and in the following, we denote $\psi(\mathbf{r}) = \psi(0, \omega)$. Two-operator expectation values are calculated with Gaussian integration of the Grassmann variables

$$\langle \psi^a \bar{\psi}^b \rangle = \int \mathcal{D}[\psi] \psi^a \bar{\psi}^b e^{i\mathcal{S}[\psi, \bar{\psi}]} = i \mathcal{G}_{ab},$$

where $a, b$ are sets of relevant indices $(i, \sigma, 1, 2)$ and $\mathcal{G}_{ab}$ denotes the matrix element of the Green’s function. For the quantum dot model, the current is given by

$$I_N = \frac{i\tau}{4} \sum_{\sigma=1,\downarrow} \int \frac{d\omega}{2\pi} \left( \langle \hat{\psi}_{Lσ}^\dagger \hat{d}_{σ} \rangle - \langle \hat{\psi}_{Lσ} \hat{d}_{σ}^\dagger \rangle - \langle \hat{d}_{σ} \hat{\psi}_{Rσ}^\dagger \rangle + \langle \hat{d}_{σ}^\dagger \hat{\psi}_{Rσ} \rangle \right),$$

and the particle density at the quantum dot by

$$\langle n_{dσ} \rangle = \frac{1}{2} \sum_{\sigma=1,\downarrow} \int \frac{d\omega}{2\pi} \left( \langle \hat{d}_{σ} \hat{d}_{σ}^\dagger \rangle - \langle \hat{d}_{σ}^\dagger \hat{d}_{σ} \rangle + \langle \hat{d}_{σ}^\dagger \hat{d}_{σ}^\dagger \rangle \right).$$
The action $S = \int \frac{d\tau}{2\pi} \bar{\Psi}(\omega) G^{-1}(\omega) \Psi(\omega)$ can be written as the sum \cite{10}
\[ S = S_L + S_R + S_d + S_I + S_{\text{loss}}, \] (8)
where the first four terms arise from the Hamiltonian and the last one from the loss term in Eq. (1). For uncoupled reservoirs, $S_{L,R}$ is written in terms of the four-component Nambu-Keldysh spinor $\Psi_i = (\psi_i^{\uparrow}, \bar{\psi}_i^{\uparrow}, \psi_i^{\downarrow}, \bar{\psi}_i^{\downarrow})^T$ and inverse Green’s function $G_i^{-1}$ has the structure
\[ G_i^{-1} = \begin{pmatrix} 0 & [g_i^{A}]^{-1} & 0 & [g_i^{R}]^{-1} \end{pmatrix}, \] (9)
where $[G_i^{A}]^K = -[g_i^{R}]^{-1} \bar{g}_i^{K} [g_i^{A}]^{-1}$, and $g_i^A$, $g_i^R$, and $g_i^K$ are the advanced, retarded, and Keldysh components of the Green’s function. For conventional $s$-wave Fermi superfluids, $g_i^A$, $g_i^R$, and $g_i^K$ are expressed by matrices of size $2 \times 2$. The functions $[g_i^{A}]^{-1}$ and $[g_i^{R}]^{-1}$ are obtained locally at $r = 0$ as \cite{11, 12}
\[ [g_i^{A,R}]^{-1} = \frac{W}{\sqrt{\Delta^2 - (\omega + \eta)^2}} \begin{pmatrix} \omega \pm i \eta & -\Delta & \omega \pm i \eta \\ \Delta & 0 & -\Delta \\ -\Delta & 0 & \Delta \end{pmatrix}, \] (10)
where $\eta > 0$ is an infinitesimal constant which regularizes the Green’s functions. The upper (lower) sign corresponds to the retarded (advanced) Green’s function. We define the frequency relative to the chemical potential as $\omega = \omega - \mu_i$. By assuming the reservoirs are in a thermal state at all times, the Keldysh component $g^K$ can be obtained through the fluctuation-dissipation theorem, $g^K = (g^R - g^A)[1 - 2n_F(\bar{\omega})]$. Here, $n_F(\omega) = (e^{\omega/T} + 1)^{-1}$ with temperature $T$ denotes the Fermi-Dirac distribution. The inverse Green’s function for the noninteracting quantum dot contains the dissipation term of Eq. (1) (see Refs. \cite{30, 31, 32}). It has the same form as Eq. (9) with the different components given by
\[ [g_i^{A,R}]^{-1}(\omega) = \begin{pmatrix} \omega - \epsilon_d \pm \frac{\gamma_d}{2} & 0 \\ -\omega - \epsilon_d \pm \frac{\gamma_d}{2} & 0 \end{pmatrix}, \] (11)
\[ [G_i^{\sigma}]^{-1}(\omega) = \begin{pmatrix} i\gamma_0 & 0 \\ 0 & -i\gamma_0 \end{pmatrix}. \] (12)

In the absence of a bias, the action including the left and right reservoirs, tunneling, and loss could be represented by an $8 \times 8$ matrix $G^{-1}$. One must take into account that in the presence of $\Delta$, fermions in either reservoir at equal but opposite frequencies $\pm \bar{\omega}$ relative to the chemical potential are coupled, while the tunneling term couples fermions on the left and right sides at the same absolute frequency $\omega$. Therefore, a finite bias leads to an action which is not block diagonal in frequency, and is represented by an infinite-size matrix with elements at increasing discrete frequencies. This corresponds physically to multiple Andreev reflections. As the off-diagonal matrix elements $[g^{A,R}]_{12}^{\pm}$, $[g^{A,R}]_{21}^{\pm}$ decay with $\bar{\omega}$, the matrix can be truncated at a certain order of the tunneling process according to desired accuracy. It can then be inverted numerically to obtain the matrix elements $G_{ab}$.

A spin bias $\Delta b = (\Delta \mu_1 - \Delta \mu_2)/2$ between the reservoirs may be modeled by including a magnetic field term in the local Green’s functions \cite{33}, so that Eq. (10) is replaced by
\[ [g^{R,A}]^{-1} = \begin{pmatrix} \frac{W(\omega + \epsilon_d)}{\sqrt{\Delta^2 - (\omega + \epsilon_d)^2}} & \frac{W \Delta}{\sqrt{\Delta^2 - (\omega - \epsilon_d)^2}} \\ \frac{W \Delta}{\sqrt{\Delta^2 - (\omega + \epsilon_d)^2}} & \frac{W(\omega - \epsilon_d)}{\sqrt{\Delta^2 - (\omega - \epsilon_d)^2}} \end{pmatrix}. \] (13)
Here, we have defined $\mu_i = (\mu_{1i} + \mu_{2i})/2$ and $b_i = (\mu_{1i} - \mu_{2i})/2$. This minimal approach corresponds to pairing with an energy offset from the Fermi level on either side, and it does not take into account the spatially modulated order parameter which typically arises from FFLO pairing at a finite center-of-mass momentum.

The current is
\[ I_N = \frac{i\tau}{2} \sum_{\sigma = \uparrow, \downarrow} \int \frac{d\omega}{2\pi} \left( \langle \psi_R^\sigma \bar{\psi}_L^\sigma \rangle - \langle \bar{\psi}_R^\sigma \psi_L^\sigma \rangle \right), \]
and the particle density in either reservoir is given by
\[ \langle n_{i\sigma} \rangle = \frac{1}{\pi} \sum_{\sigma = \uparrow, \downarrow} \int \frac{d\omega}{2\pi} \left( \langle \psi_{i\sigma} \bar{\psi}_{i\sigma} \rangle - \langle \bar{\psi}_{i\sigma} \psi_{i\sigma} \rangle + \langle \psi_{i\sigma}^2 \bar{\psi}_{i\sigma}^2 \rangle \right). \]

For numerical calculation of such observables, the frequency integral must be truncated to a range $\bar{\omega} \in [-\Lambda, \Lambda]$ given by a high energy cutoff $\Lambda$. The integrand of the current $I_N$ falls off quickly for $|\bar{\omega}| > \Lambda$ so $I_N$ is independent of $\Lambda$. The density $\langle n_{i\sigma} \rangle$ on the other hand diverges with $\Lambda$ as a consequence of the linear dispersion assumed here, so $\Lambda$ is treated as a fitting parameter to agree with the measured atom loss rate $N$. Note that, in the quantum dot model, the occupation of the quantum dot is bounded due to its Lorentzian lineshape, though it contains contributions from an unphysically large range of energies since the fitted linewidth is several times the gap. Therefore the cutoff $\Lambda$ is also needed in the quantum dot model.

II. EXPERIMENTAL DETAILS

A. Experimental cycle

We begin by preparing a non-degenerate, spin-balanced cloud with approximately $2 \times 10^5$ atoms in the first and third-lowest spin states of $^6\text{Li}$ (|$\uparrow$⟩ and |$\downarrow$⟩ respectively) at their Feshbach resonance $B = 689.7$ G. The cloud is trapped along $y$ (direction of transport) by a magnetic trap and along $x$ and $z$ by an optical dipole trap propagating along $y$. We then prepare the initial imbalance $\Delta N$ by shifting the cloud along $y$ with a magnetic field gradient, then ramping up the power of the beams that define the QPC along with a repulsive “wall” beam to block transport (Sec. II B), and finally returning the cloud center to coincide with the QPC. We then evaporatively cool the cloud to degeneracy by ramping down the dipole trap power with a magnetic field gradient along gravity to further lower the trap depth. At
this stage, just before transport, the trap frequencies are $\nu_{r,\text{trap}} = 184.7(6)$ Hz, $\nu_{y,\text{trap}} = 28.2(1)$ Hz, $\nu_{z,\text{trap}} = 178.4(5)$ Hz and the cloud has $N = 195(14) \times 10^3$ atoms at $T = 100(2) \, \text{nK}$, corresponding to $T/T_F = 0.256(8)$, and an atom number imbalance of $\Delta N/N = 0.373(4)$. There is a very small spin imbalance $\Delta M = M_L - M_R \lesssim 0.02N$ that is due to a small overall imbalance in the populations of the two spin states $M = N_L - N_T \lesssim 0.02N$.

Because we prepare the density imbalance before evaporation, the evaporation efficiencies of the two reservoirs differ due to the differing atom number in each one. This leads to a large initial temperature bias between the two reservoirs $\Delta T$ which gives rise to a strong, unwanted thermoelectric effect [44]. To suppress this effect, we apply a magnetic field gradient during evaporation to compress the reservoir with lower atom number and decompress the one with more atoms, thereby equalizing their evaporation efficiencies such that $\Delta T \approx 0$ at the beginning of transport.

We then ramp up the powers of the attractive gate and dissipation beams. Then transport is allowed by switching off the wall beam before switching it on at a later time to block transport again. All beam powers except the dipole trap and the wall are ramped down such that each reservoir is in a half-harmonic trap, at which point we take an absorption image of both spin states by sending two pulses of light resonant with each spin state in quick succession (225 µs apart) synchronized to a CCD camera operating in fast kinetics acquisition mode.

B. Defining the quantum point contact

Two TEM$_{01}$-like beams of repulsive 532 nm light define the QPC at the intersection of their nodal planes. The first propagates along $x$, has a Gaussian waist $w_{x,QPC} = 30.2$ µm along $y$, and has a peak confinement frequency of $\nu_{x,QPC} = 9.9(2)$ kHz along $z$. The second propagates along $z$, has a Gaussian waist of $w_{z,QPC} = 6.82$ µm along $y$, and a peak confinement frequency of $\nu_{z,QPC} = 11(2)$ kHz along $x$ for the data presented. The light intensity is not exactly zero at these beams’ nodal planes which leads to an additional contribution to their zero-point energy that is non-negligible. This is calibrated by measuring the onset of quantized conductance in a non-interacting system ($B = 568$ G) under the same conditions [45]. The wall beam is a 532 nm elliptical beam that also propagates along $z$ and has a waist of 8.58 µm along $y$.

The gate beam of attractive 777 nm light propagates along $z$ and has Gaussian waists of $w_{x,\text{gate}} = 30.3$ µm along $x$ and $w_{y,\text{gate}} = 31.7$ µm along $y$. The potential energy it exerts on the atoms $V_g < 0$ was determined by calibrating its power at the location of the atoms and computing the resulting ac-Stark shift from its known beam profile and wavelength. As this beam is essentially a weakly-focused optical tweezer, it has a transverse confinement frequency $\nu_{x,\text{gate}} = \sqrt{|V_g|/m}/\pi w_{x,\text{gate}} = 220(2)$ Hz which adds in quadrature to $\nu_{z,\text{QPC}}$ to determine the net confinement frequency at the center of the QPC. We observed that this power calibration drifted over the course of our measurements. As a result, there is a relatively large systematic uncertainty on this gate potential $V_g$ of approximately 6%. The equilibrium imbalance $\Delta N$ is extremely sensitive to the center position of this beam along $y$, especially in the presence of dissipation, on a scale of $\sim 1$ µm – many times smaller than the beam’s waist. We therefore align this beam position such that, starting from zero imbalance $\Delta N = 0$, the imbalance remains zero at long time for each dissipation strength.

C. Dissipation mechanism

The dissipative beam is a tightly-focused Gaussian beam propagating along $z$ aligned to the center of the QPC with waists $w_{x,\text{dis}} = 1.28$ µm and $w_{y,\text{dis}} = 1.31$ µm. It is shaped and positioned using a digital micromirror device whose Fourier plane is projected onto the plane of the QPC by a high-NA microscope objective [46]. Its frequency is tuned into resonance with the $\sigma^-$-transition between $|\downarrow\rangle = |1/2, 1\rangle - \epsilon |1/2, 0\rangle$ and $|e\rangle \approx |3/2, 0\rangle$ (written in the basis $|m_f, m_i\rangle$). 2.98 GHz blue-detuned from the center of the D$_2$ line [25]. The excited state $|e\rangle$ has an inverse lifetime of $\Gamma_e/2\pi = 5.87$ MHz and decays into the first $|1\rangle = |\downarrow\rangle$, fifth $|5\rangle = |1/2, 0\rangle + \epsilon |1/2, 1\rangle$, and sixth $|6\rangle = |1/2, 1\rangle$ lowest ground states with branching ratios $\Gamma_{e1} = 2.9 \times 10^{-3} \Gamma_e$, $\Gamma_{e5} = 0.997\Gamma_e$, and $\Gamma_{e6} = 8 \times 10^{-5} \Gamma_e$. This ensures that only 0.29% of $|\downarrow\rangle$ atoms that have scattered a photon return to $|\downarrow\rangle$ and the rest are optically pumped to $|5\rangle$. These states are then out of resonance, anti-trapped by the magnetic field as they are low field-seeking, and have a recoil energy larger than the optical trap depth so they are quickly lost from the system. We therefore interpret a photon scattering event by a $|\downarrow\rangle$ atom as a quantum jump of the loss process modeled in Sec. I. We use this optical pumping process instead of photon scattering on a closed transition since the strongly-interacting $|\downarrow\rangle$ atoms imparted with the photon recoil energy quickly destroy the cloud.

We calibrate the power $P_{\text{dis}}$ of this beam at the location of the atoms and use the known beam shape to compute its peak intensity $I_{\text{dis}} = 2P_{\text{dis}}/(\pi w_{x,\text{dis}}^2 w_{y,\text{dis}})$, from which we can compute the photon scattering rates $\Gamma_{\sigma}$ and ac-Stark shifts $V_{\sigma}$ of $|\downarrow\rangle$ and $|\uparrow\rangle$ [25] which are directly proportional to $I_{\text{dis}}$, since all intensities used in our measurements are less than 3% of the saturation intensity of this transition $I_{\text{sat}} = 8.7$ kW/m$^2$. In addition to the intensity and frequency, it is also crucial to know the polarization of the beam to compute the photon scattering rates and ac-Stark shifts. We measured the fraction of the total power in each polarization directly on the atoms by monitoring atom loss rate in a non-interacting system. We tuned the beam’s frequency into resonance with the $\sigma^-$, $\pi$, and $\sigma^+$ transitions of $|\downarrow\rangle$’s $m_J = -1/2$
component and measured the atom loss rate at constant power by fitting an exponential to $N_i(t)$. The ratio between each loss rate scaled by each transition’s oscillator strength is then the ratio of the powers in each polarization. This calibration yielded $78.7(5)\%$ of the total power in $\sigma^+$, $6.6(1)\%$ in $\pi$, and $14.8(6)\%$ in $\sigma^-$. Given these calibrations and a photon scattering rate of $\Gamma_i$, we have $\Gamma_i = 5.4 \times 10^{-4} \Gamma_1$, $V'_i = k_B \times (1.49 \text{mK/kHz}) \Gamma_1$, and $V_i = k_B \times (1.37 \text{mK/kHz}) \Gamma_1$ which are all negligible compared to other relevant scales.

In the non-interacting gas ($B = 568 \text{ G}$), we observe that photon scattering of $|\downarrow\rangle$ does not induce any loss of $|\uparrow\rangle$—strong evidence for the interactions between $|\uparrow\rangle = |3\rangle$ and $|\downarrow\rangle = |5\rangle$ being weak as expected [47], however at unitarity ($B = 689.7 \text{ G}$), there are significant losses of $|\uparrow\rangle$ and its loss rate is proportional to the loss rate of $|\downarrow\rangle N_\uparrow \approx 0.7 N_\downarrow$. This indicates that the loss is due to the $s$-wave interactions between $|\downarrow\rangle$ and $|\uparrow\rangle$. This is supported by the observation that the relative loss ratio $r = N_\uparrow/N_\downarrow$ decreases with increasing temperature and decreasing chemical potential where interatomic binding is weaker. Microscopically, a photon absorption by $|\downarrow\rangle$ is a Rabi $\pi$ pulse to the excited state $|e\rangle$. Concurrently with this $\pi$ pulse, the photon momentum is transferred to the atom and the interaction strength ($s$-wave scattering length) is ramped down due to the internal state of the atom changing. This means that fraction of the photon recoil is transferred to $|\uparrow\rangle$ and the energy of the system is changed by sweeping the $s$-wave scattering length, so, according to Tau’s sweep theorem, the energy transferred to the system is proportional to the contact [48]. At the level of our theoretical model, this means that $|\uparrow\rangle$ experiences a “mean-field dissipation” with rate $\gamma_1 r \gamma_1 (N_\downarrow)$.

### III. RESERVOIR THERMODYNAMICS

For each shot, we obtain the column density $n_{i\sigma}^{\text{col}}(y, z)$ of both half-harmonic reservoirs $i = L, R$ and both spin states $\sigma = \downarrow, \uparrow$ from an absorption image taken in situ along the $x$-direction with a calibrated imaging system. The atom number in each spin state and reservoir is determined by integrating the density over the half-plane of each reservoir

$$N_{i\sigma} = \int_0^\infty dy \int_{-\infty}^{\infty} dz n_{i\sigma}^{\text{col}}(y, z).$$

We then compute the second spatial moment of the line density

$$\langle y^2 \rangle_{i\sigma} = \frac{\int_{-\infty}^{\infty} dy n_{i\sigma}^{\text{lin}}(y) y^2}{\int_{-\infty}^{\infty} dy n_{i\sigma}^{\text{lin}}(y)}$$

which determines the total energy per particle through the virial theorem for the spin-balanced, harmonically-trapped unitary fermi gas $E_{i\sigma}/N_{i\sigma} = 3m \omega_y^2 \langle y^2 \rangle_{i\sigma}$ [49]. With the $N_{i\sigma}$ and $E_{i\sigma}$ along with the average trap frequency $\bar{\omega} = (\omega_x^2 \omega_y^2 \omega_z^2)^{1/3}$, we use the equation of state (EoS) to solve for the temperature $T_{i\sigma}$.

Since the dissipation rate is spin-dependent, a spin polarization builds up over time and the virial theorem is no longer valid. However, since the magnetization in each reservoir is small $M/N < 10\%$ for most of the data, the shape of the minority spin cloud do not change significantly [50, 51] and we can apply the same temperature extraction procedure with limited systematic error. This approach is validated by the observation that the extracted temperature does not depend on dissipation strength or transport time within statistical measurement fluctuations. This observation also indicates that thermoelectric effects are negligible in our parameter regime (e.g. density or spin currents coupling to entropy currents that lead to a temperature bias).

With the measured atom number in both spin states and their temperature in each reservoir $N_{i\downarrow}, N_{i\uparrow}, T_i$, we use the recently-computed EoS of the spin-polarized gas at finite temperature [52] to solve for the chemical potential of both spins in the reservoir $\mu_{i\sigma}$. This determine the average chemical potential $\mu_L = (\mu_{\downarrow} + \mu_{\uparrow})/2$ and the Zeeman field $b_L = (\mu_{\downarrow} - \mu_{\uparrow})/2$ which, together with the temperature $T_i$, completely characterize the pressure equation of state of the spatially homogeneous system

$$P(\mu, b, T) = \frac{1}{\beta \lambda_i^2} f_P(\beta \mu, \beta b)$$

where $\lambda_T = \sqrt{2\pi \hbar^2/m k_B T}$ is the thermal de Broglie wavelength and $\beta = 1/k_B T$ is the inverse temperature. The free energy of a harmonically-trapped system $\Omega(\mu, b, T)$ can be computed from $P(\mu, b, T)$ (i.e. “free energy density”) by applying the local density approximation with spatially-varying average chemical potential

$$\mu(r) = \mu - \frac{m}{2} \left( \omega_x^2 x^2 + \omega_y^2 y^2 + \omega_z^2 z^2 \right)$$

and spatially constant $b$ and $T$, yielding

$$-\Omega(\mu, b, T) = \int d^3r P[\mu(r), b, T]$$

$$= \left( \frac{k_B T}{\hbar \omega} \right)^4 \frac{2}{\sqrt{\pi}} \int_0^\infty d\eta \sqrt{\eta} f_P(\beta \mu - \eta, \beta b)$$

$$= \frac{(k_B T)^4}{(\hbar \omega)^4} f_\Omega(\beta \mu, \beta b).$$

From $f_\Omega$, we can compute the atom number $N = N_\downarrow + N_\uparrow$ and magnetization $M = N_\uparrow - N_\downarrow$ in the harmonic trap

$$N(\mu, b, T) = -\frac{\partial \Omega}{\partial \mu} = \left( \frac{k_B T}{\hbar \omega} \right)^3 \frac{\partial f_\Omega}{\partial (\beta \mu)}$$

$$M(\mu, b, T) = -\frac{\partial \Omega}{\partial b} = \left( \frac{k_B T}{\hbar \omega} \right)^3 \frac{\partial f_\Omega}{\partial (\beta b)}$$

and thermodynamic response functions like the compressibility $\kappa$, spin susceptibility $\chi$, and “magnetic-
pressibility” α

\begin{align*}
\kappa(\mu, b, T) &= \frac{\partial N}{\partial \mu} = \frac{(k_B T)^2}{(\hbar \omega)^3} \frac{\partial^2 f_\Omega}{\partial (\beta \mu)^2} \\
\chi(\mu, b, T) &= \frac{\partial M}{\partial b} = \frac{(k_B T)^2}{(\hbar \omega)^3} \frac{\partial^2 f_\Omega}{\partial (3b)^2} \\
\alpha(\mu, b, T) &= \frac{\partial N}{\partial b} = \frac{\partial M}{\partial \mu} = \frac{(k_B T)^2}{(\hbar \omega)^3} \frac{\partial^2 f_\Omega}{\partial (\beta \mu) \partial (3b)}
\end{align*}

which set the relationship between the atom number and magnetization imbalances and the chemical potential and Zeeman field biases in the linear response regime of the reservoirs

\[
\left( \frac{\Delta N}{\Delta M} \right) \approx \frac{1}{2} \left( \frac{\kappa}{\alpha} \chi \right) \left( \frac{\Delta \mu}{\Delta b} \right).
\]

A misalignment of the center of the magnetic trap with respect to the QPC leads to a finite atom number imbalance at equilibrium \(\Delta N(N_\mu = 0) = \Delta N_\infty \neq 0\). This can be compensated in the equation of state \(\Delta \mu(\Delta N, N, T)\) by \(\Delta \mu(\Delta N - \Delta N_\infty, N, T)\) where \(\Delta N_\infty\) is fixed to \(\Delta N\) at the longest transport time.

We find that all of our data lies within the computed range of the Zeeman field \(|\beta b| < 2\) (Sec. IV C). Below the minimum computed chemical potential \(\beta \mu < -5\), we use the third order virial expansion for \(f_p(\beta \mu, |\beta b|)\) [52]. The densest regions of the cloud however lie above the critical degeneracy of \(\beta \mu \approx 2.5\) [53] where the computed EoS is known to deviate from the true response of the system. We therefore take the EoS to be that of the low temperature superfluid with phononic excitations [54]

\[
f_p^{SF}(\beta \mu, |\beta b|) = \frac{16}{15 \sqrt{\pi}} \left[ \frac{(\beta \mu)^{5/2}}{\xi(\beta \mu)^{5/2}} + \frac{\pi^4}{96} \frac{3}{(\beta \mu)^{3/2}} \right]
\]

where \(\xi = 0.370\) is the Bertsch parameter [55]. The finite-temperature superfluid is known to have finite spin susceptibility [56, 57] so using this EoS, which has vanishing spin susceptibility \(\partial^2 f_p^{SF}/\partial b^2 = 0\), introduces some systematic error. However, the susceptibility in this regime is exponentially suppressed by a large pairing gap so the overall error in \(f_{\Omega}\) is negligible.

As the phase boundary between the normal and superfluid phases is not exactly known, we take it to be the intersection of \(f_p\) and \(f_p^{SF}\) such that \(f_p\) is continuous but its first derivative is discontinuous. This choice agrees well with a recent functional renormalization group approach to compute the phase boundary [58]. This choice also minimizes the potentially detrimental artifacts that could arise in the thermodynamic response functions which are derivatives of \(f_{\Omega}\). In fact, there is no discontinuity in \(\kappa, \chi\), or \(\alpha\), likely because \(f_{\Omega}\) is a weighted integral of \(f_p\) and therefore naturally smoother.

### IV. DATA ANALYSIS

#### A. Estimate of the number of transport modes

An important parameter that contributes to the overall timescale of transport is the number of available transport modes in the QPC \(n_m\) [59]. If the modes are decoupled from each other and the current they each carry is identical, then the total current is \(n_m\) times the current of a single mode \(-1/2\Delta \omega^K = n_m I_N^K\). In the non-interacting system, where it is possible to exactly compute the current in the Landauer formalism with our detailed knowledge of the potential energy landscape in the QPC [25], \(n_m\) naturally emerges as the sum of the Fermi-Dirac occupation of each transverse mode indexed by the harmonic numbers \(n_x, n_z\)

\[
n_m = \sum_{n_x, n_z} \frac{1}{1 + \exp \left\{ (E_{n_x, n_z}(0) - \mu) / k_B T \right\}}
\]

where \(E_{n_x, n_z}(0)\) is the energy of mode \(n_x, n_z\) at the center of the QPC. This expression is valid if the mode transmits particles adiabatically (de-coupled modes) [60] which we have verified from our measurements on the non-interacting gas. The energy has contributions from the confinement along \(x\) and \(z\) from the two QPC beams and the gate beam as well as the residual repulsive light on the QPC that increases the energy. These two contributions are proportional to the square root of the power of the beam and linear in the beam power respectively. Finally including the gate potential \(V_g\) from the attractive beam, the energy can be written

\[
E_{n_x, n_z}(0) = \sqrt{(a_x \sqrt{P_x})^2 + \nu_{x, gate}^2 (n_x + 1/2)} + b_x P_x + a_z \sqrt{P_z (n_z + 1/2)} + b_z P_z + V_g
\]

where \(P_x, P_z\) is the power of the QPC beam that confines along the \(x, z\) direction and the parameters \(a_x, b_x, a_z, b_z\) were calibrated by measuring \(E_0(0)\) in a non-interacting system (the onset of quantized conductance) with various \(P_x\) and \(P_z\) and fitting Eq. 24. Note that this estimation method assumes the same effective tunneling amplitude \(\tau\) of each mode and neglects the possibility that the zero-point energy of the modes is reduced by the strong inter-particle attraction (cf. Ref.[61]).

#### B. Estimate of the superfluid gap

In solid state systems, where the superconducting reservoirs and contacts are geometrically well-defined, the gap \(\Delta\) that enters the theoretical model is given by the bulk of the reservoir material [62]. However, in our system where the separation between reservoirs, contact, and channel are not sharp but vary smoothly over the length scales of the Gaussian beams that define the QPC, this choice is not as obvious.
We therefore estimate \( \Delta \) from the local chemical potential at the most degenerate location in the system at the contacts of the 1D region \( \mu_c = \max_r [\mu - V(r)] \) where \( V(r) \) includes all traps, zero-point energies \( E_{0,0}(r) \) of each confinement beam, and the gate potential \( V_g(r) \). \( \mu_c \) then gives the local density \( n(\mu_c, T) = \partial P/\partial \mu \) from the equation of state of the homogeneous 3D system, and the Fermi energy \( E_F = \hbar^2(3\pi^2n)^2/2m \approx \mu_c/\xi \) determines the gap \( \Delta = 0.44E_F \) [63]. The temperature \( k_B T/E_F < 0.04 \) is low enough to safely apply the zero temperature value of \( \Delta \). Note that this estimation method neglects the possibility that the additional confinement in the 2- and 1-dimensional regions enhances the pairing gap [61].

C. Fitting procedure

For a given data set (the time evolution of the density imbalance with a fixed configuration of the QPC, gate, and dissipation strength), we first fit the time evolution of the total atom number \( N(t) = N_L(t) + N_R(t) \). As shown in the main text, the model cannot quantitatively predict the loss rates \( \dot{N}_c \) so we use a phenomenological model inspired by our previous work on a non-interacting gas under similar conditions [25]. There, the loss rate \( \dot{N} \) was found to be proportional to the fraction of atoms with energies above the zero point energy of the channel, which itself is approximately proportional to the atom number \( N \). Moreover, we can assume that a \(|\uparrow\rangle\) atom is lost with some probability \( r \) for each dissipated \(|\downarrow\rangle\) atom, motivated by the dissipation mechanism described in Sec. II C where some fraction of the photon recoil is transferred from \(|\downarrow\rangle\) to its paired \(|\uparrow\rangle\) via interactions. The parameter \( r \) can therefore be interpreted as a measure of the spin correlations at the location of the dissipation. These two assumptions lead to the rate equations

\[
\frac{d}{dt} N(t) = -\frac{1}{\tau_N} N(t)
\]

\[
\frac{d}{dt} N_i(t) = r \frac{d}{dt} N_i(t)
\]

(25)

with fit parameters \( \tau_N \) and \( r \) (the loss ratio) as well as the initial atom number in each spin state. We only fit the data before \( \Delta N \) fully equilibrates (\( \Delta N/N \geq 0.02 \), on the order of our experimental stability of \( \Delta N/N \)) to obtain a better fit result at initial times, in which we are most interested. These fits for the data presented in Fig. 2 in the main text are shown in Fig. 5(a,b). We see that even this simple model fits the data well, allowing us to reliably extract \( \dot{N} \) and \( \dot{M} \) from the data.

Now with the fitted \( N(t) \) and \( M(t) \), we turn to fitting \( \Delta N(t)/N(t) \) and \( \Delta M(t)/N(t) \) (we fit the relative imbalances \( \Delta N/N \) and \( \Delta M/N \) rather than the absolute imbalances \( \Delta N \) and \( \Delta M \) as this normalizes out the shot-to-shot atom number fluctuations of \( \sim 10\% \)). Using the EoS for the half-harmonic reservoirs described in Sec. III, the measured atom number \( N_i(t) (i = L, R) \), magnetization \( M_i(t) \), and temperature \( T_i(t) \) in each reservoir at a given time \( t \) can be converted to a chemical potential \( \mu_i(t) \) and Zeeman field \( b_i(t) \). The biases between the reservoirs \( \Delta \mu = \mu_L - \mu_R \) and \( \Delta b = b_L - b_R \) and the average temperature \( T = (T_L + T_R)/2 \), along with the superfluid gap \( \Delta \) and number of modes \( n_m \) (both determined by the average chemical potential \( \mu = (\mu_L + \mu_R)/2 \) and the gate potential correction factor \( \eta_g \)), tunneling amplitude \( \tau \), and dissipation rates \( \gamma_L \) and \( \gamma_R \) are then the inputs to our model to compute the density and spin currents \( \dot{I}_N^K \) and \( \dot{I}_M^K \) (Sec. I). In this way, the decay of the density and magnetization imbalances is described by the system of

![FIG. 5. Atom number (a), magnetization (b) and relative magnetization (c) over time for varying dissipation strength. Data sets in (a) and (b) are vertically shifted for clarity by +10 and -0.02 units respectively where the \( \Gamma_1 = 0 \) data is the most shifted in (a) and unshifted in (b). The extracted fit parameters defined in Eq. 25 and 27 are plotted in the insets. Inset of (a) is identical to the normalized atom loss rate plotted in Fig 4(a). The horizontal bar in the inset of (b) represents the weighted average \( r \) used in the fit of \( \Delta N/N \).](image-url)
non-linear differential equations

\[
\begin{align*}
- \frac{1}{2} \frac{d}{dt} \Delta N^K & = n_m I_N^K (\Delta N^K, \Delta M^K, N, M, T, \eta_g, \tau, \gamma_1, \gamma_2) \\
- \frac{1}{2} \frac{d}{dt} \Delta M^K & = n_m I_M^K (\Delta N^K, \Delta M^K, N, M, T, \eta_g, \tau, \gamma_1, \gamma_2)
\end{align*}
\]

subject to the initial conditions \(\Delta N^K(0) = \Delta N(0)\) and \(\Delta M^K(0) = \Delta M(0)\), which can be easily solved numerically. However, we find that for values of \(\gamma\) that reproduce the observed \(I_N\), the computed \(I_N^K\) underestimates the observed \(I_M = -\Delta M/2\) by over an order of magnitude. This means that we cannot obtain good fits by simultaneously fitting \(\Delta N(t)/N(t)\) and \(\Delta M(t)/N(t)\).

The evolution of the relative magnetization \(\Delta M/N\) is shown in Fig. 5(c) along with a fit to a phenomenological rate equation assuming linear spin transport

\[
\frac{d}{dt} \left( \frac{\Delta M(t)}{N(t)} \right) = -\frac{1}{\tau_{\Delta M}} \left[ \frac{\Delta M(t)}{N(t)} - \left( \frac{\Delta M}{N} \right)_{\infty} \right].
\]

Similarly to how a chemical potential bias \(\Delta \mu\) can drive a magnetization current \(I_M\) due to the spin asymmetry \(\gamma_1 - \gamma_\gamma\), the Zeeman field bias induced by the magnetization imbalance \(\Delta b \approx \Delta M/\chi\) can drive a particle current \(I_N\). Furthermore, \(\Delta M\) can directly influence \(\Delta \mu\) via the “magnetic compressibility” \(\Delta \mu = \Delta N/\kappa - (\alpha/\kappa \chi) \Delta M\) (Sec. III). Because the model cannot accurately reproduce the off-diagonal conductance \(\partial I_M/\partial \Delta \mu\), we cannot expect it to compute \(\partial I_N/\partial \Delta b\) either. However, on general grounds, we can expect that \(\partial I_N/\partial \Delta b \approx \partial I_M/\partial \Delta \mu\) – essentially a statement of Onsager’s reciprocal relations for coupled density and spin transport. Our data shows that, for the strongest dissipation, \(\partial I_M/\partial \Delta \mu \sim -n_m/\hbar\) and therefore the contribution of \(\Delta b\) to \(I_N\) is \(\leq -n_m \Delta b/\hbar\). The \(\Delta b\) extracted from \(\Delta M\) shown in Fig. 5(c) is \(\Delta b/\Delta < 0.03\) and is far below this limit for all but the longest transport times at the strongest dissipation, so the influence on \(I_N\) is below even the normal part of the chemical-potential driven transport \(\sim n_m \Delta \mu/\hbar\), let alone the superfluid part \(\sim n_m \Delta b/\hbar\).

Furthermore, the influence on \(\Delta \mu\) driven by \(\Delta M\) is \((\alpha/\kappa \chi) \Delta M < 0.005 \Delta\) in the most extreme case which is negligible compared to the density imbalance-induced bias \(\Delta N/\kappa \sim 0.05 \Delta\). In summary, the influence of the magnetization dynamics \(\Delta M(t)\) on the density dynamics \(\Delta N(t)\) is negligible to a good approximation.

Therefore, because \(\Delta N(t)\) is de-coupled from \(\Delta M(t)\) and \(I_N^K\) is insensitive to \(\gamma_1 - \gamma_\gamma\), we fit only \(\Delta N(t)/N(t)\) with our model for \(\Delta N^K(t)\) using the spin-balanced equation of state \((b = \Delta b = 0)\) and with the ratio of the dissipation strengths fixed to the fitted loss ratio \(\gamma_1/\gamma_\gamma = r = \hat{N}_I/\hat{N}_t\)

\[
- \frac{1}{2} \frac{d}{dt} \Delta N^K = n_m I_N^K (\Delta N^K, N, T, r, \eta_g, \tau, \gamma_1, \gamma_\gamma) - \frac{1}{2} \frac{d}{dt} \Delta M^K = n_m I_M^K (\Delta N^K, N, M, T, \eta_g, \tau, \gamma_1, \gamma_\gamma).
\]

As explained in the main text, we first fit the imbalance decay curves where the dissipation beam was not present such that we can fix \(\gamma_1 = \gamma_\gamma = 0\). The fit parameters here are the average tunneling amplitude of the occupied modes \(\tau\) and, due to its experimental uncertainty (Sec. II B), the correction factor \(\eta_g\) of the gate potential \(V_g\) which enters into \(n_m\) and \(\Delta\). The fitted values of \(V_g\) are consistent with its calibrated value for the data shown and within 20\% for other measurements with different QPC parameters (\(\nu_g\) and \(V_g\)), supporting this approach. Although the fitted \(r\) shows a small dependence on \(\Gamma_1\), we simply use the average value for the fit of \(\Delta N/N\) due to the uncertainty in the fit of \(r\) and the fact that the calculated current is insensitive to \(r\). With \(r, \tau\) and \(\eta_g\) fixed, we then fit \(\gamma_1\) (which determines \(\gamma_\gamma = \tau r\gamma_1\)) to each data set taken under the same QPC and gate conditions but with varying dissipation strength.

D. Numerical derivation of current

To compare the data with the theory in current-bias relation [Fig. 2(b)], we obtain the current \(I_N\) from a numerical derivative of the measured \(\Delta N\). In practice, we pass the relative imbalance data \(\Delta N(t)/N(t)\) [Fig. 2(a)] into a Savitzky–Golay filter of order 3, from which the first derivative at each time point is given by the 4 neighbouring points. \(\Delta N\) is then obtained assuming the fitted exponential decay of \(N\). To estimate the uncertainty of the numerical time derivative, we apply a bootstrap-type method. We pass an averaged time evolution with a random sampling without replacement of 70\% of the entire dataset to the Savitzky–Golay filter to have an ensemble of derivative results (100 realizations). We then take the standard deviation of this ensemble to estimate its uncertainty. There is however still a visible oscillation in the derivative data beyond the estimated uncertainty, which is a typical artifact due to the limited sampling in time in the original data.

E. Removing outliers

Due to some technical issues such as atom number fluctuation and a sporadic synchronization failure of our spatial light modulator that generates the dissipative beam, we have data points that are clearly beyond statistical uncertainties. We remove the outliers beyond 3\(\sigma\) of the statistical distribution of \(\Delta N/N\) of the entire dataset in a given QPC and dissipation configuration, and those beyond 3\(\sigma\) in \(N, N_j/N_t\), and \(\Delta M/N\) which are better indicators of technical problems. As we have different transport time in the dataset, we use a smoothing spline to determine the mean value of the data as function of the transport time. This procedure assumes no knowledge of our theoretical model.