Time-resolved Coulomb collision of single electrons

A series of recent experiments have shown that collision of ballistic electrons in semiconductors can be used to probe the indistinguishability of single-electron wavepackets. Perhaps surprisingly, their Coulomb interaction has not been seen due to screening. Here we show Coulomb-dominated collision of high-energy single electrons in counter-propagating ballistic edge states, probed by measuring partition statistics while adjusting the collision timing. Although some experimental data suggest antibunching behaviour, we show that this is not due to quantum statistics but to strong repulsive Coulomb interactions. This prevents the wavepacket overlap needed for fermionic exchange statistics but suggests new ways to utilize Coulomb interactions: microscopically isolated and time-resolved interactions between ballistic electrons can enable the use of the Coulomb interaction for high-speed sensing or gate operations on flying electron qubits.

Precise control over interactions between ballistic electrons will enable us to exploit Coulomb interactions in novel ways, to develop high-speed sensing, to reach a nonlinear regime in electron quantum optics and to realise schemes for fundamental two-qubit operations on flying electrons. Time-resolved collisions between electrons have been used to probe the indistinguishability, Wigner function and decoherence of single-electron wavepackets. Due to the effects of screening, none of these experiments were performed in a regime where Coulomb interactions are particularly strong. Here we explore the Coulomb collision of two high-energy electrons in counter-propagating ballistic edge states. We show that, in this kind of unscreened device, the partitioning probabilities at different electron arrival times and barrier heights are shaped by Coulomb repulsion between the electrons. This suggests a new class of devices for studying and manipulating interactions of ballistic single electrons.

In principle, time-resolved electronic interactions can be studied with a wavepacket collider sketched in Fig. 1. Single-electron sources S1 and S2 emit particles with relative delay Δt into an experimentally defined collision region. Quantum exchange statistics can affect how particles are partitioned into detectors D1 and D2 for different injection times. This Hong–Ou–Mandel collision geometry originates in quantum optics where it was first used to probe photon wavepackets via quantum exchange statistics. When photons collide at a beam splitter they display a bunching behaviour related to their bosonic nature. In the electronic equivalent, fermionic exchange effects can create an antibunching of wavepackets, detected via reduced current noise at the detectors. This effect can be used as a measurement of the indistinguishability of the wavepackets, an important figure of merit for quantum coherent transport. However, in general, understanding the behaviour of the electrons in the collision region is not straightforward if Coulomb effects are present in addition to exchange statistics.

For sources injecting electrons near the Fermi energy, the impact of the Coulomb interaction on the electron trajectory is diminished by screening. Changes to the single-electron trajectories or velocity by Coulomb interactions between single electrons has not been seen, although coupling to nearby conducting channels has been detected via decoherence of the wavepackets. Where electronic density is

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isolated from the time-averaged current. We study the partitioning giving interactions. For each source separately by varying the barrier height of each source. The emission times from the two sources can be fine-tuned so that the nominal difference between arrival times at the centre barrier Δt. The emission times from the two sources can be fine-tuned so that the barrier height is Δt = 0. Additionally, at delay times near Δt = 0 the maximum partition noise S12/Δt is reduced from the independent scattering limit by up to Δt ∼ 30%, as shown in Fig. 3c. This effect is driven by an increase in the number of events where both electrons are reflected due to a repulsive interaction. Although the noise reduction superficially resembles a partial Pauli dip, we show below that the full partitioning statistics, calculated from the time-averaged current and noise data, reveal a Coulomb-dominated system.

A clear piece of evidence for the dominance of the Coulomb interaction is that interaction effects persist for grossly mismatched injection energy. Figure 3d shows that the individual transmission thresholds for each source are associated with E1 and E2 (ref. 7) are perturbed by interactions when arrival is synchronized. The size and polarity of this effect is consistent with the shift in the peak of the noise data when E1 = E2 as in Fig. 3b,c, suggesting that they are both caused by Coulomb repulsion. As injection at grossly different energies makes electrons distinguishable, only a manifestation of the long-range Coulomb interaction can explain these effects.

To build a more detailed picture of this interaction, we use the measured noise S1 and current imbalance ΔI, to compute the partition statistics, the probability of various charge distributions in the detector contacts. The partition probabilities for the three possible detector charge states P0, P1, P2, where Ij is the number of electrons in detector D1,D2. A particular set of Pij values uniquely corresponds to a certain current noise and time-averaged current and is a powerful parameterization of the scattering outcomes. For overlapping trajectories, the transmission probability T from a specified source is not accessible. In contrast, the partition statistics Pij captures where charges terminate and not from which source they came, so these can always be calculated. As P1 + P20 + P02 = 1 only two values, for example, P02 and P20, are required to fully capture the partitioning outcomes.

Partition statistics are shown in Fig. 3e–j where each panel has different time delay Δt. Values of (P02,P20) are plotted parametrically; each point is taken at a different barrier height as the barrier is lowered in the direction indicated by the arrows. The representative non-interacting cases in Fig. 3e (Δt = 0 ps) have P02,P20 = 0, P11 = 1 for high and low barriers and the expected statistical mixture of outcomes when E1 ≈ E2, namely P02,max = P20/min = 0.25 and P02,min = 0.5. The results in the interacting regime (Δt ∼ 10 ps) are shown in Fig. 3f–j. Preferential transmission into D1 or D2 caused by interactions causes an imbalance in P02 and P20 (a bias toward the P20 or P02 axis). P20,max and

is shown in Fig. 3a. For values near Δt = 0 we see a pronounced current imbalance with ΔI ≈ 0 for Δt > 0 and ΔI ≈ −0.2 for Δt < 0. The polarity of this signal corresponds to preferential transmission of the earlier electron.

Measuring the time-dependent noise current, ∆C1,∆C2, using the additional connections shown in Fig. 2b enables us to measure correlations in partitioning outcomes from individual scattering cycles that are not accessible in a time-averaged measurement. The cross-correlated noise current S12, simultaneously measured in contacts D1 and D2 (Methods), is mapped as a function of E1 and Δt in Fig. 3b. When Δt ≫ 0, electrons arrive at very different times and do not interact. The non-interacting noise is explained by independent stochastic transmission of electrons from each source given by S12 = 2e2∫ (T1(T1 - T2)) (ref. 8). This gives S12 ≈ 0 when the barrier is either much higher or lower than the injection energy, but a peak value S12/Δt = e2f when E1 = E2. We find that this signature is modified by Coulomb interactions between the electrons when their arrival is synchronized.

For values of Δt = 0 (that is, electrons collide) the position of the noise peak moves to a lower value of the programmed barrier height E1. This apparent shift in detected energy, E1 → E1 and E2 → E2, is a manifestation of the Coulomb repulsion between the electrons which gives a transient increase in the effective barrier height of −0.5–1 meV. Conservation of energy is preserved by a reduction in the kinetic energy corresponding to the increased Coulomb repulsion. Additionally, at delay times near Δt = 0 the maximum partition noise S12/Δt is reduced from the independent scattering limit by up to ∼30%, as shown in Fig. 3c. This effect is driven by an increase in the number of events where both electrons are reflected due to a repulsive interaction. Although the noise reduction superficially resembles a partial Pauli dip, we show below that the full partitioning statistics, calculated from the time-averaged current and noise data, reveal a Coulomb-dominated system.

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Fig. 2 | Electron collider implemented with high-energy single electrons. a. Partitioning of electrons in a high-energy device. Three outcomes are characterized by charge distributions (2,0) when both electrons enter D1, (0,2) when both enter D2, and (1,1) when one enters each. The proximity of electrons is exaggerated for clarity (see Methods for timing information).

b. False colour scanning electron microscopy image of overall device structure and measurement system. See Methods for details. Electrons emitted from sources S1 (left, red) and S2 (right, blue) follow trajectories indicated by arrows, colliding at the central barrier (green). The dashed region indicates the collision region shown in a. Inset: energy level drawing to show adjustable ejection energy.

Fig. 3 | Current imbalance, current noise suppression and partition statistics. a. Current imbalance $\Delta I_{21}$ for different barrier heights $E_B$ and time delays $\Delta t$. Favoured events associated with the negative and positive current imbalances are indicated schematically. b. Map of current noise $S_{12}$ for barrier height $E_B$ and time delay $\Delta t$. Dashed line is a guide to peak noise position. Colour-coded arrows indicate the positions of data shown in panels e–j. c, S$_{12}$ in the non-interacting ($\Delta t = -20$ ps) and strongly interacting regime ($\Delta t = +2$ ps). Errors bars are the s.e.m. from $n > 11$ measurements (1 min each) at each voltage in multiple sweeps (Supplementary Information, section 1.2). The reduction in $S_{12}$ is associated with an increase in repulsion events as indicated. d. Current measurements $P_{20}^C$ and $P_{12}^C$ for varying barrier height for grossly mismatched energy difference $E_1 - E_2 = 5$ meV. e–j. Each panel shows partition probabilities ($P_{20}, P_{12}$) at different barrier heights for a chosen $\Delta t$ (see arrows in b): +20 ps (e), −20 ps (f), −9 ps (g), 0 ps (h), +2 ps (i), +9 ps (j). Error bars are derived from the standard error in the measured current and noise. Underlying colour map shows noise $S_{12}$ (contour lines are at 0.1e$^2$/intervals). e shows non-interacting cases (electrons miss each other) at $\Delta t = -20$ ps (dark grey) and +20 ps (light grey). Arrows show the sweep from high to low barrier height, reaching the maximum noise at the independent scattering limit (thick solid line). At this point ($P_{20}, P_{12}$) = (0.25,0.25) (circled) there is an equal statistical mixture of each outcome event. f–j, as e but for selected interacting cases at values of $\Delta t$ corresponding to arrows in panel b.
We show below that the partitioning statistics capture a detailed signature of the Coulomb interaction in this system. The full partitioning dataset can be modelled by considering microscopic trajectories which follow the $E \times B$ drift motion. The shape of these trajectories is set by the partitioning barrier potential and the Coulomb interaction, which is only weakly screened. We also show that the small timing and energy fluctuations of the source are important in determining the statistical outcomes seen experimentally.

We compute the trajectories for an assumed form of the potential $U_2(x, y)$ near the barrier and an interaction potential $U_{int}(\vec{r}_1, \vec{r}_2) \propto \beta/\vec{r}$ at short range ($\beta$ is the strength of the repulsion), which but weakens gradually at longer distances due to screening by the metallic surface gates located at a distance of ~90 nm above the two-dimensional electron gas (2DEG) (Methods)\(^2\). Coulomb interactions are diminished for distances much larger than this. This limits the interaction to the region near the effective tunnelling point, enabling us to take a saddle point barrier potential as an approximation. Calculations based on the $E \times B$ trajectories are valid in the regime of our experiments where the energy window of quantum scattering (over which the barrier transmission probability changes from 0 to 1) is much smaller than the energy uncertainty of the electrons and/or the purity of the electrons is sufficiently low (Supplementary Information, section 2).

We estimate the Coulomb interaction strength by comparing our theory (based on the trajectories and the assumed interaction form in equation (4) of Supplementary Information, section 2.2) and the experimental data at $E_1 \neq E_2$ (as in Fig. 3d); the detailed estimation can be found elsewhere. The estimated strength is consistent with the bare value of $\beta = \epsilon_2/(4\pi \epsilon_0 \epsilon_r) = 112$ meV nm, where $\epsilon_0$ is the vacuum permittivity and $\epsilon_r = 12.9$ is the relative permittivity of GaAs. Below we show good agreement between the calculation and experimental data using this value, supporting our conclusion that the Coulomb interaction is the dominant factor.

Example trajectories are shown in Fig. 4 for injection energy exceeding the barrier height $E_1 = E_2 = 0.5$ meV. This illustrates three important cases: completely mismatched arrival times in Fig. 4a; near-synchronized arrival in Fig. 4b,c; and exactly synchronized arrival in Fig. 4d. In non-interacting cases such as Fig. 4a, trajectories simply follow the equipotential contours of $U_2$ and are transmitted or reflected depending on injection energy alone. In Fig. 4a both are transmitted, giving $\Delta t_j = 0$. At or very near perfect synchronization as in Fig. 4d, both electrons are reflected as they undergo a Coulomb collision centred at the beam splitter, also giving $\Delta t_j = 0$. The distance of closest approach, typically 20–50 nm for the parameters used here, is governed by the injection energy, the saddle potential and the Coulomb interaction.

For the near-synchronized cases in Fig. 4b,c interactions modify the trajectory of the electrons in a more complicated way in that the outcome depends upon the polarity of the relative arrival time $\Delta t$, that is, which electron reaches the barrier first. The late-arriving electron (from Si in these examples) tends to be deflected by the electron which enters the barrier region first (here from S2)\(^2\). This asymmetric effect leads to $\Delta t = 1$ for $\Delta t > 0$ and $\Delta t = -1$ for $\Delta t < 0$ for a narrow range of $\Delta t$, as in Fig. 4b,c. Detecting this order-of-arrival effect would demonstrate a control of the Coulomb repulsion on time scales relevant for transient interactions with device components in edge states. To explore this possibility, we numerically compute a large number of trajectories at different injection energies $E_1 - E_0$, $E_2 - E_0$ and relative arrival time $\Delta t$. This enables us to make a detailed map of the charge partitioning outcomes (Supplementary Information, section 2) and use this to reveal how details of the Coulomb interaction appear in experimental partitioning data.

The current imbalance $\Delta I_j$ is plotted in Fig. 4c, computed for trajectories with different arrival time delays and barrier heights for equal injection energy $E_j = E_n$. The parameters for the specific example trajectories in Fig. 4a–d is indicated with markers. This map consists of four regions whose boundaries are defined by the values of $E_{th} = E_{th}'$ the programmed barrier height at which the first or second electrons are no longer transmitted in the presence of Coulomb repulsion. The functional form of $E_j(\Delta t)$ is an asymmetric peak created by the accumulated effect of the Coulomb interaction during the interaction. In the region where the peaks overlap, there is a regime where the electrons repel most strongly, as in Fig. 4d. At larger values of $|\Delta t|$, when electron arrival is not exactly synchronized, the asymmetry gives $E_j(\Delta t) \neq E_n(\Delta t)$. This is because the effective kinetic energy loss of each particle is different. This creates an intermediate region where only the late-arriving electron is blocked. This predicts a current imbalance in $I_j$ with the same polarity as that measured experimentally in Fig. 3a. This suggests that the fine structure in the time-resolved Coulomb collision described above is indeed visible experimentally.

In this calculation, we find that the exact values of $E_j$ and the range $\Delta t$ of $\Delta t$ over which the interaction extends are governed by the cyclotron frequency, the curvature of the saddle point $\omega_{xy} = \sqrt{\omega_x \omega_y}$ and the value of the screening length scale. However, the qualitative behaviour of the partitioning map is not sensitive to the precise parameters used. In the experiment, extrinsic emission time and energy fluctuations of the source broaden the boundaries of the different scattering outcomes as the barrier height and time of arrival vary. This has an impact on the maximum visibility of the order-of-arrival effects, completely smearing them out in the case of grossly time-broadened emission (Supplementary Information, section 1.7). We compute the partition probabilities $P_{21}, P_{11}, P_{00}$ from the expected ensemble average of the scattering outcomes using initial conditions $(E_1, E_2, \Delta t)$ of the trajectories obtained from the source energy–time distribution (that is, energy, time widths $\sigma_e$ and $\sigma_t$ measured for each source; Methods). For the device shown here $\sigma_e = 1.7$ ps, $\sigma_t = 5.2$ ps. This is larger than our estimate of the intrinsic size of the interaction feature $\tau_e < 2$ ps and therefore this broadening is an important effect.

The full experimental partition statistics and the computed ensemble average are shown together in Fig. 4f–h. The agreement between the experimental data (upper panels) and model calculations (lower panels) show that the Coulomb interaction, broadened by the fluctuations of the electron sources, is accessible experimentally. Overall, the ensemble effects blur the sharp boundaries of the classical outcomes, reducing the size of the shift in the noise peak. The energy–time correlation present in the emission distributions, a characteristic feature of the source\(^1\), and mismatches between the exact source emission distributions explain the fine structure of the partitioning data. The energy shift and the order-of-arrival effect we see show that the sharpness of our emission distributions are on the same picosecond time scale required to control single-electron interactions with high fidelity.

Single-electron antibunching signatures driven by the Pauli effect were observed with mesoscopic capacitor sources which alternately emit electrons and holes\(^2\). In that case, as expected of fermionic identical-particle exchange statistics, noise suppression was seen for electron–electron synchronization but not for electron–hole synchronization\(^1\). This was an important observation to rule out the possible roles of Coulomb effects. We have to consider here the reverse question: in a system with relatively strong Coulomb interactions, is it ever possible for fermionic exchange effects to play a role? The technical factors that might prevent this include the limited purity of the sources\(^1\) and the energy selectivity of the barrier, which modifies the electronic wavepacket\(^11,12\). More fundamentally, it appears from our measurements and a realistic model that the trajectories of the electrons are so strongly repelled that achieving a high overlap of the incoming wavepackets is unlikely, at least in this parameter regime, and there would be no manifestation of fermionic exchange effects even for perfect sources and an energy-independent barrier.
As far as we have seen, Coulomb effects are visible provided the fluctuations in arrival difference $\Delta t$ of the sources are not too much larger than the interaction time scale $\tau$ of the barrier. The collision technique actually provides a very sensitive readout of the emission-time distribution of single-electron sources, an important figure of merit for on-demand electron sources, without the bandwidth limitations of gate modulation techniques. This approach can be used to optimize the emission characteristics of on-demand sources for applications in time-resolved sensing and other high-speed applications. Measurements with strongly mismatched energy as in Fig. 3d may also enable a mapping of the two-dimensional potential landscape (modelled by the saddle point curvature) of the barrier. This possibility arises because the interacting trajectories are sensitive to the shape of potential barrier in a way that depends on the exact injection energies, which can be tuned.

In summary, we have revealed the effect of mutual repulsion between single electrons in high-energy chiral ballistic channels using an electron collider. In previous work, performed in a highly screened regime, this approach has been used to probe indistinguishability of identical quantum particles. Here, experiments in a depleted channel at high magnetic fields reveal Coulomb-dominated perturbations of electron trajectories. It will be interesting to further trace the perturbation of the trajectories by performing tomographic measurements. At much lower energy, the interplay of the fermionic exchange statistics and the Coulomb interactions may be accessible in a weak interaction regime achieved by electrostatic gate arrangements.

We note here that we have recently learned of experiments performed with electrons sources similar to this, but operating slightly closer to the Fermi energy (30–60 meV) using single-shot charge detection (N. Ubbelohde et al., unpublished observations). These experiments also suggest a Coulomb-dominated regime that extends over a broad range of energy. Experiments with electrons confined in travelling surface acoustic wave minima (J. Wang et al., unpublished observations) also shown an apparent antibunching that may be driven by Coulomb repulsion. This shows that the Coulomb interaction may dominate in a range of scenarios relevant for electron quantum optics or flying qubits.

**Online content**

Any methods, additional references, Nature Portfolio reporting summaries, source data, extended data, supplementary information, acknowledgements, peer review information; details of author contributions and competing interests; and statements of data and code availability are available at https://doi.org/10.1038/s41565-023-01369-4.

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Methods
Device
Our sources are fabricated on GaAs/AlGaAs with a 2DEG wafer depth $d_{2DEG} = 90$ nm, carrier density $n_{2D} = 2.4 \times 10^{11}$ cm$^{-2}$, mobility $\mu = 2.5 \times 10^{4}$ cm$^2$/Vs. Each pump is a gate-defined quantum dot in a channel 800 nm wide at a distance of 5 µm from the centre barrier. Entrance gates, adjacent to the source reservoir, are controlled by $V_{\text{SLG}}(t)$, modulated to load electrons from the source reservoirs. The same modulation pushes the trapped electron towards the exit barriers controlled by $V_a$ and $V_d$ (refs. 7,25,26). Interaction effects have been explored with both sinoidal and tailored pump waveforms at several frequencies. Data presented here are from a sinuouslidal drive at $f = 500$ MHz. The device is measured in a perpendicular magnetic field $B = 10$ T (into the plane of Fig. 2) in a dry dilution refrigerator with mixing chamber $T_{\text{m}} \leq 0.1$ K.

Energy tuning and measurement
The pumping exit barrier controls the number of electrons collected and the energy of the injected electrons. This is used to adjust/match the injection energies of the two sources. Relative electron energy can be measured from the setting of $E_a$ required to block transmission. Gate voltage values scaled to energy using the phonon emission features in the energy distribution at multiples of $\hbar \omega = 36$ meV (ref. 19). The absolute injection energy requires separate calibration experiments not performed here, but the presence of $n_{2D} \simeq 3$ side-bands to the main electron distribution at a spacing of $\hbar \omega_a$ indicates an energy of at least $\sim 100$ meV as in previous reports.

Edge state transport
In a perpendicular magnetic field, $B = 10$ T, the electrons emitted from the pump follow a chiral edge channel towards the centre of the device. Electrons follow high-energy Landau levels with drift velocity $v_d = (E \times B)/B^2$ governed by the edge electric field $E$ and magnetic field $B$. Relaxation to lower-energy states is strongly inhibited by the edge gates, here controlled by $V_{\text{NV}}, V_{\text{V}}$ and $V_{\text{edge}}$ as described in Supplementary Information, section 1.4. A small residual phonon emission means that -5% of the injected electrons are always reflected. This creates a small residual noise when $E_a - \hbar \omega_a < E_b < E_a$ visible as a residual value of $P_{\text{c}}$ as $P_{\alpha} > 0$ in Fig. 3e–f for low barriers. This effect is included in the computed partition statistics in Fig. 4.

Time-of-arrival synchronization
To find the delay setting corresponding to $\Delta t = 0$ a synchronization pulse is applied to the centre barrier synchronized with both electron pumps but with an adjustable delay $t_d$. Maps of drain current versus $E_a$ and $t_d$ reveal the shape of the synchronization pulse, referenced to each electron’s arrival time (see ref. 1 and Supplementary Information, section 1.6). Note that the transit time estimated from the velocity $t_{\text{path}} \leq 100$ ps is shorter than the repeat period $1/f_d = 2$ ns, and hence there are no interactions between electrons emitted in different cycles.

Partition readout
Average current in output channels $I_{\alpha 1}$ and $I_{\alpha 2}$ is read with current-to-voltage converters. Using $E_{\text{c}} = E_{\text{c}} + E_{\text{e}} = 2E_{\text{f}}$, the data in Fig. 3a are calculated from $\Delta I_{\alpha} = \left(E_{\alpha 2} - E_{\alpha 1}\right)/2E_{\text{f}}$, where $E_{\alpha 1}$ and $E_{\alpha 2}$ are measured on D1 and D2, calibrated using single-electron partitioning. Here we consider the absolute value only; the measured value of $\Delta I_{\alpha}$ is negative due to the detector configuration. See Supplementary Information, sections 1.2 and 1.3 for details.

Calculation of trajectories
The trajectories of electrons in a 2D plane are obtained based on a total potential $U_{\alpha}(x, t) = U_{\alpha 0}(x) + U_{\alpha 1}(x, t) + U_{\alpha 2}(x, t)$ with $U_{\alpha 0} = E_a - m^* \omega_a x^2/2 + m^* \omega_y^2 y^2/2$ Numerical integration of the equations of motion gives trajectories and the final charge distribution can be evaluated after electrons leave the scattering area (Supplementary Information, section 2). These calculations are similar to those of ref. 13. We define $U_{\alpha}$ to include a contribution from the interaction between electrons in the GaAs material ($c_\ell = 12.9$) and an interaction with an image charge at a distance $2d_{2DEG}$ from the plane (ref. 27) which screens the interaction at distances $r_{\text{sc}} > d_{2DEG} \sim 90$ nm. This weakens the interaction and makes the effects of the potential away from the scattering region less relevant. Selection of initial particle positions sets both electronic energy and the relative injection time (this uses the analytical solutions to the non-interacting equations of motion appropriate for large $r_{\text{sc}}$), see Supplementary Information, section 2. We choose a single value of $\omega_a = \omega_a = 5 \times 10^4$ s$^{-1}$ to reproduce the experimentally observed features after accounting for the effect of the input energy–time distributions (see below). The qualitative features of the scattering outcomes map (the regions labelled in Fig. 4e) do not change appreciably in this 2D regime with exact choice of $\omega_a, \omega_b$ (ref. 13). The temporal width of the dip in $E_1$ and $E_2$ become narrower with increasing $\omega_a$, as the electrons are forced closer together.

Ensemble average over input energy, time distributions
To simulate the effect of stochastic fluctuations in the initial energy–time distribution we first compute classical outcomes for discrete values of different input energies $E_1, E_2$ and times $t_1, t_2$ in a high-resolution grid, then sample these outcomes using distributions based on experimentally observed input distributions for the two sources $W(E_1, t_1)$ and $W(E_2, t_2)$ (Supplementary Information, section 2). The partition noise and current imbalance are readily computed from the sampled outcomes, from which the partition statistics can be calculated. Rather than a full tomography of the incoming Wigner distributions we use a simpler technique which parameterizes the Wigner distribution as a bivariate Gaussian $W_{\alpha}(\sigma_1, \sigma_2)$ with energy and time widths $\sigma_1$ and $\sigma_2$ and energy–time correlation coefficient $r$.

$$W_{\alpha} = \frac{1}{2\pi\sigma_1\sigma_2} \exp \left[ -\frac{1}{2(1-r^2)} \left( \frac{E^2}{\sigma_1^2} + \frac{t^2}{\sigma_2^2} - \frac{2Ert}{\sigma_1\sigma_2} \right) \right]$$

We find the size of the emission distributions similar to that found previously with the most important difference being the temporal length of the two electrons is different, $\sigma_1 = 1.7 \pm 0.5$ ps and $\sigma_2 = 5.2 \pm 1$ ps, while the injected energy distributions are similar, $\sigma_{E_1} = 0.85 \pm 0.2$ meV and $\sigma_{E_2} = 1.05 \pm 0.2$ meV (Supplementary Information, section 1.7). The correlation coefficients are $r_1 = 0.5$ and $r_2 = 0.85$. We compute the partition statistics over $N = 10^4$ samples from distributions with these parameters over the maps of scattering outcomes.

Data availability
The experimental data that support the findings of this study are available in the SEQOUIA community repository at https://zenodo.org/communities/sequoia/at https://doi.org/10.5281/zenodo.7643880. Calculation details are provided in Supplementary Information.

Code availability
The code used for analyses and figures is available from the corresponding author upon reasonable request.

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Author contributions
J.D.F. designed and developed the measurement system and experimental methodology; with input from P.S. and M.K., designed samples; performed experiment, data acquisition and analysis; with W.P., helped to explore numerical calculation of particle trajectories and comparison with data. W.P. developed a method to calculate classical particle trajectories; performed numerical calculations to compare with experimental data; investigated methods to establish the validity of the classical model. S.R. developed a one-dimensional quantum collision model (published separately) and helped to develop a classical method of calculation of particle trajectories and advised on its realm of validity. P.S. developed fabrication techniques for electron pump devices; and fabricated samples used in this paper and preliminary batches of prototype electron colliders. J.P.G. and G.A.C.J. provided electron beam sample patterning. I.F. and D.A.R. provided wafers for the device substrates. H.-S.S. oversaw the development of both quantum and classical models of particle collision. M.K. directed this research project; supported J.D.F. for experiments and data analysis, and suggested the underpinning mechanisms that result in positive and negative correlation in two-electron transmission/reflection observed experimentally. The manuscript was written by J.D.F., M.K., W.P., S.R. and H.-S.S. with review by other authors.

Competing interests
The authors declare no competing interests.

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