

## Loss and Decoherence at the Quantum Hall-Superconductor Interface

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We perform a systematic study of Andreev conversion at the interface between a superconductor and graphene in the quantum Hall (QH) regime. We find that the probability of Andreev conversion from electrons to holes follows an unexpected but clear trend: the dependencies on temperature and magnetic field are nearly decoupled. We discuss these trends and the role of the superconducting vortices, whose normal cores could both absorb and dephase the individual electrons in a QH edge. Our Letter may pave the road to engineering a future generation of hybrid devices for exploiting superconductivity proximity in chiral channels.

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Combining superconductors and quantum Hall (QH) systems has been proposed as a particularly promising direction for creating novel topological states and excitations [1]. Over the past few years, significant progress has been achieved in developing such hybrid structures [2–14]. In particular, hybridization of QH edge states across a narrow superconducting wire is expected to create a gapped topological superconductor [6,15]. In the fractional QH systems, the strong interactions potentially fractionalize Majorana fermions into parafermions [13,16], a key ingredient for universal topological quantum computing [1] and exotic circuit elements such as fractional charge transistors [15].

At the interface between a superconductor and a QH system, the QH edge states are expected to be proximitized, turning into chiral Andreev edge states (CAES). These are dispersive states which hybridize the electron and the hole amplitudes [17]. An electron approaching the superconducting region is converted to a linear combination of CAES, which interfere as they propagate along the interface. The outgoing particle can either stay as an electron or turn into a hole. We have previously observed clear evidence of the electron-hole conversion in the quantum Hall devices with superconducting contacts [12]. However, the exact mechanism of this conversion remains open: the role of the disorder, superconducting vortices, and the exact nature of the CAES in this system have been discussed. It is known that in order to observe a strong Andreev conversion in an ideal system a precise matching between the superconductor and the QH edge state momenta is required [18–21]. However, the presence of disorder is expected to relax this constraint [22,23]. Vortices in the superconductor and extra QH channels induced by doping could also modify the signal [22–25]. In particular, the normal cores of

the vortices can absorb the electrons and holes, reducing the amplitude of the measured signal [12,23,26].

To address the microscopic mechanisms affecting the Andreev conversion, here we perform a systematic study of the conversion probability vs temperature  $T$ , magnetic field  $B$ , and interfacial length  $L$ . We find that the dependence of the electron-hole conversion probability on magnetic field and temperature nearly factorizes. We suggest a simple phenomenological expression involving exponential decays as a function of  $B$ ,  $T$ , and  $L$ , and a prefactor determined by the configuration of superconducting vortices. The expression captures the observed dependencies very well, and can be interpreted in terms of the CAES loss and decoherence. We finally discuss the distribution of the Andreev conversion probability [23], which is unexpectedly found to have a roughly triangular shape. We discuss the implications of our findings for the future development of the more complex devices, which will further explore the physics of superconducting correlations in the chiral states.

The main device studied here is a hBN/graphene/hBN heterostructure in contact with both superconducting and normal electrodes [Fig. 1(a)]. The superconducting electrode (light gray) is made of sputtered Mo-Re alloy (50-50 in weight) with a critical temperature  $T_c \sim 10$  K and an upper critical field  $H_{c2}$  exceeding 12 T. The work function of the alloy  $\sim 4.2$  eV [27] is slightly lower than that of graphene  $\sim 4.5$  eV [28], resulting in an  $n$ -doped graphene region nearby. The graphene sheet is separated into two independent regions by the etching step that defines the superconducting electrode. The widths of the resulting superconductor-graphene interfaces are 0.5 and 1  $\mu\text{m}$ . The normal contacts (yellow) are thermally evaporated Cr/Au. The electron density inside the graphene is controlled by applying a voltage  $V_G$  to the graphite gate that spans the

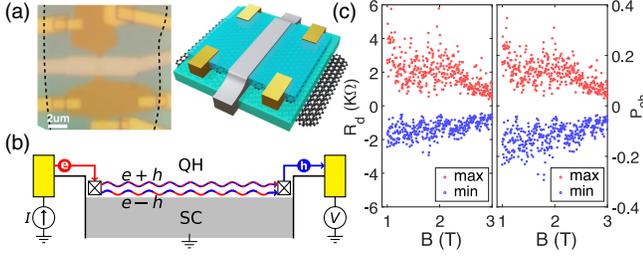


FIG. 1. (a) Optical image (left) and 3D schematics (right) of the device. The black dashed lines label the boundaries of the graphite gate underneath. The yellow electrodes are normal Cr/Au contacts and the light gray is the superconducting MoRe, which forms 0.5 and 1  $\mu\text{m}$  interfaces with graphene. (b) Sketch of the measurement setup. The current is injected from the upstream contact into the grounded superconductor, while measuring the voltage at the downstream contact. (c) The maximum and minimum values of  $R_d$  (left) and  $P_{eh}$  (right) measured for the shorter interface. The statistical information is collected in the  $V_G$  range corresponding to  $\nu = 2$ , as outlined by the green lines in Fig. 3(a). The temperature is 40 mK.

whole area underneath the heterostructure. While such graphite gates are known to efficiently screen the disorder potential, the results here are similar to our previous measurements of samples without the graphite gate. This indicates that the observed physics is not strongly influenced by the disorder in the graphene layer.

We measure the nonlocal resistance downstream of a grounded superconducting contact,  $R_d = dV_d/dI$ , as sketched in Fig. 1(b) and demonstrated in the Supplemental Material [29]. The CAES formed by the superconductor travel along the interface and recombine into either an electron or hole (or their linear combination) at the end of the interface. Sweeping the gate voltage  $V_G$  on top of the QH plateau tunes the momentum difference between the interfering CAES and produces an oscillating pattern of  $R_d(V_G)$  (Fig. S1 in [29]) [12]. Because of the disordered nature of the interface, these interference patterns are highly irregular and resemble the universal conductance fluctuations [22,23]. Throughout the Letter we analyze the statistical properties of these patterns.

We further convert  $R_d$  into the difference between the probabilities of normal and Andreev reflections,  $P_{eh} \equiv P_e - P_h$ , where  $P_e$  ( $P_h$ ) is the probability of an electron (or a hole) to be emitted downstream of the superconductor. It is straightforward to show that  $P_{eh} = R_d/(R_d + R_H)$ , where  $R_H$  is the Hall resistance [12]. Note that the extreme values of  $R_d$  are reached either for the pure electron reflection ( $P_e = 1$ ,  $P_h = 0$  and  $R_d = \infty$ , the interface is effectively fully opaque), or for perfect Andreev conversion ( $P_e = 0$ ,  $P_h = 1$ ,  $R_d = -R_H/2$ , a Cooper pair is transferred across the interface per incoming electron). We therefore expect that the distribution of  $R_d$  should be skewed toward positive resistances.

Indeed, this skewness can be observed by studying an imbalance between the maximum and minimum values of

the downstream resistance  $R_d(V_G)$ . We extract these quantities in the  $V_G$  range corresponding to the  $\nu = 2$  plateau for a given field and then plot  $R_{\text{max}}$  and  $R_{\text{min}}$  as a function of  $B$  in Fig. 1(c). The field changes the vortex configuration every few mT [12], thereby allowing us to sample multiple different patterns of  $R_d(V_G)$ . The data clearly show that  $R_{\text{max}}$  is on average larger than  $|R_{\text{min}}|$ . The apparent imbalance between electrons and holes is eliminated by converting  $R_d$  to  $P_{eh}$  in Fig. 1(c). (See Fig. S3 in [29] for a similar result measured in another device.) We conclude that the probabilities of an electron or a hole being emitted downstream ( $P_e$  and  $P_h$ ) are very similar. From now on, we present  $P_{eh}$  in lieu of  $R_d$ .

Depending on the length of the interface, the highest electron-hole conversion efficiency we observe is about 0.2–0.3. While being very high, this value is still far from reaching unity. Thermal smearing, decoherence and tunneling into the normal vortex cores can all contribute to this suppression. Since at zero temperature only the effect of the vortices remains, we first look into the dependence of  $P_{eh}$  on temperature. Figure 2(a) plots  $P_{eh}$  of the shorter interface (0.5  $\mu\text{m}$ ) measured over the  $\nu = 2$  plateau at  $B = 2.2$  T from 40 to 670 mK. As the temperature  $T$  increases,  $P_{eh}$  gradually decays towards zero, except around 400 mK where the traces briefly jump to a completely different oscillation pattern. We attribute this jump to a temporary change of the configuration of superconducting vortices during the measurement.

To get a quantitative understanding of the thermal effects, we plot the standard deviation of the traces  $\sigma$  as a function of temperature in Fig. 2(b) for both interfaces at various magnetic fields. The solid triangles (open circles)

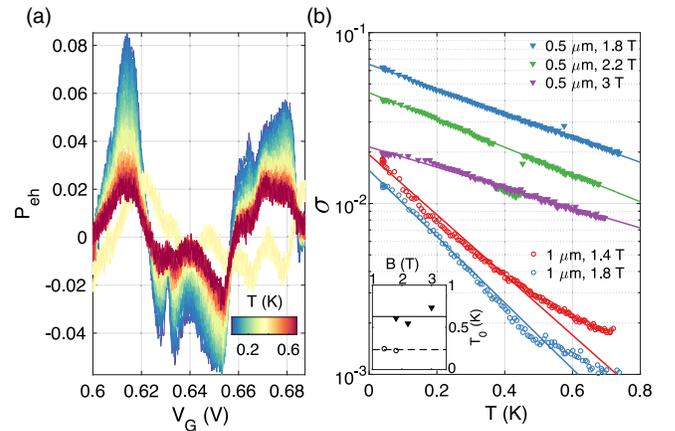


FIG. 2. (a)  $P_{eh}$  measured for the 0.5  $\mu\text{m}$  interface as a function of  $V_G$  on top of the  $\nu = 2$  plateau at  $B = 2.2$  T. The temperature is varied from the base temperature of 40 to about 670 mK. (b) The standard deviation of  $P_{eh}$  of the 0.5  $\mu\text{m}$  (filled triangle) and 1  $\mu\text{m}$  (open circle) interfaces as a function of temperature at various magnetic fields. The inset plots the decay constant  $T_0$  obtained from exponential fits vs the magnetic field. The lines represent the averages of the dots.

represent  $L = 0.5 \mu\text{m}$  ( $1 \mu\text{m}$ ). The green triangles correspond to the data in Fig. 2(a). We find an exponential decay of  $\sigma$  as a function of temperature, i.e.,  $\sigma \propto \exp(-T/T_0)$ , where  $T_0$  is the decay constant. Remarkably,  $\sigma$  follows roughly the same decay rate for a given interface length. Even the curves of Fig. 2(a) which experienced a random jump follow the same slope. (Note the few green symbols nearly overlapping with the purple ones around 0.4 K.) This means that the configuration of superconducting vortices does not have a strong influence on the temperature dependence, even if it dramatically affects the amplitude and the pattern of fluctuations.

The exponential decay is observed regardless of the length of the interface and the magnetic field. For the longer interface, the decay rate becomes less steep at higher temperatures, but only as the signals approach the noise floor of a few  $\times 10^{-3}$ , so we extract the slope from the low-temperature range. In the inset of Fig. 2(b), we plot the resulting  $T_0$  vs  $B$  for  $L = 0.5$  and  $1 \mu\text{m}$ . It is clear that  $T_0$  does not show any strong dependence on  $B$  and scales inversely with  $L$ .

In principle, the exponential temperature dependence could originate from the averaging over energy, like the thermal smearing of oscillations in a QH Fabry-Perot interferometer. However, the effects of thermal smearing are expected to saturate below  $T_0$ , while we observe an exponential decay down to temperatures more than 10 times lower (for the shorter interface). Moreover, thermal smearing is predicted to be rather inefficient in this system [23], because Andreev conversion is only weakly energy dependent. Furthermore, thermal broadening would smear the fluctuations of  $P_{eh}$  over  $V_G$ . Instead, we observe that the fluctuations in Fig. 2(a) uniformly decay with temperature with no noticeable smearing. In fact, we show that curves in Fig. 2(a) could be rescaled to nearly match (see Fig. S5 in [29]).

We therefore ascribe the temperature decay of  $P_{eh}$  to the decoherence of the CAES. Indeed, decoherence resulting in the exponential suppression of the oscillations  $\propto \exp(-T/T_0)$  has been measured in Mach-Zehnder interferometers [31,32]. However, the coherence length of the QH interferometers usually exceeds the length of the superconducting contact by more than an order of magnitude, including our own measurements using similar graphene samples [33]. The rapid decoherence of the CAES as compared to the QH edge states is likely due to their hybrid nature: The electron and the hole components should acquire an opposite phase in the fluctuating electrostatic potential. The source of these fluctuations could be the gate, or the vortex cores, which serve as a reservoir of normal electrons located next to the CAES.

Since  $T_0$  is not strongly affected by the magnetic field and the jumps of the vortex configuration, the effects of the vortices should be primarily encoded in the zero

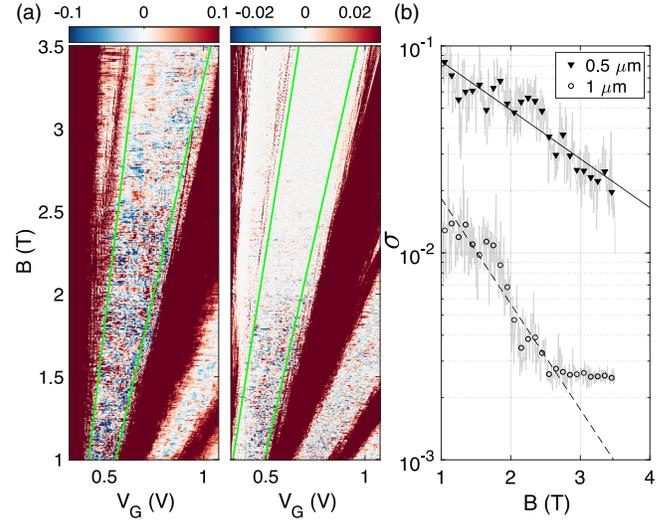


FIG. 3. (a)  $P_{eh}$  plotted against  $B$  and  $V_G$  for  $L = 0.5 \mu\text{m}$  (left) and  $1 \mu\text{m}$  (right) at  $T = 40$  mK. The green lines mark the region of the  $\nu = 2$  plateau used for the analysis in panel (b) and Fig. 1(c). (b) The standard deviation  $\sigma$  of  $P_{eh}$  measured over the  $V_G$  range corresponding to  $\nu = 2$ . The top and bottom gray curves correspond to the  $L = 0.5$  and  $1 \mu\text{m}$  interfaces. The lower curve saturates upon reaching the noise floor of the experiment, at  $\sigma \approx 2.5 \times 10^{-3}$ . The symbols (circles and triangles) represent  $\sigma$  averaged over 100 mT range,  $\langle \sigma \rangle_{\Delta B}$ . This averaging reveals the overall decay trend of  $\sigma$ , represented by the exponential fits (straight solid and dashed lines).

temperature value of  $\sigma$ . The base temperature of  $40 \text{ mK} \ll T_0$  allows us to use the corresponding values of  $\sigma$  as a close approximation to zero temperature value. In the following section, we study the dependence of this quantity on the magnetic field.

In Fig. 3(a) we show the patterns of  $P_{eh}$  over a wide range of magnetic field for the  $0.5 \mu\text{m}$  interface (left) and the  $1 \mu\text{m}$  interface (right). A corresponding Hall conductance map is shown in Fig. S2 [29]. The boundaries of the  $\nu = 2$  plateau are labeled by the green lines. The amplitude of  $P_{eh}$  for the longer interface is not only smaller but also decays much faster with increasing  $B$ . This is consistent with our picture of electrons and holes tunneling from the edge state into the normal cores of the superconducting vortices. Indeed, the probability grows with both  $L$  and the vortex density, which increases with  $B$ .

While the dependence of  $P_{eh}$  on  $B$  is highly stochastic [Fig. 1(c), and S6 in [29]], we can analyze the large-scale trend by averaging  $\sigma$  over a relatively small range of magnetic fields,  $\Delta B = 100$  mT. The resulting quantity averages over multiple vortex configurations and will be denoted by  $\langle \sigma \rangle_{\Delta B}$ . As shown in Fig. 3(b), it roughly follows an exponential decay with  $B$  for both interfaces. Excluding the points that reach the noise floor, an exponential fit produces  $B_0 = 1.84$  T for  $L = 0.5 \mu\text{m}$  and  $B_0 = 0.85$  T for  $L = 1 \mu\text{m}$ , in agreement with the expected relation  $B_0 \propto 1/L$ .

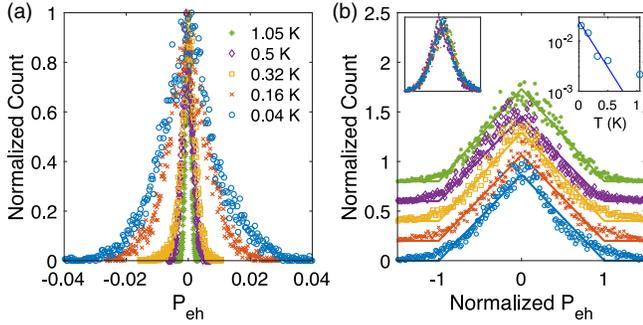


FIG. 4. (a) Histograms of  $P_{eh}$  collected for the longer interface in the 1.4–1.5 T window on the  $\nu = 2$  plateau at various temperatures. (b) The same data as in (a) rescaled by fitting a phenomenological triangular shape (lines) and normalizing  $x$  axis by the fitted width of the triangle. The curves are offset by 0.2 along the  $y$  axis. Left inset: the same normalized distributions replotted without an offset. Right inset: the fitted width of the triangle plotted vs  $T$ . The line is a guide for the eye, which follows the decay constant  $T_0$  obtained in Fig. 2(b).

Note that this relation further implies an exponential decay of  $\langle \sigma \rangle_{\Delta B}$  as a function of  $L$ ,  $\langle \sigma \rangle_{\Delta B} \propto \exp(-L/L_0(B))$  with  $L_0 \propto 1/B$  at zero temperature. This dependence indicates that the electrons/holes are efficiently absorbed by multiple vortices, so that the inverse decay length  $1/L_0$  is proportional to the vortex density  $\propto B$  [23]. The  $\exp(-L/L_0)$  dependence is directly confirmed in another device with three contacts of different interface length, as shown in Fig. S4 [29]. Interestingly, both devices have similar  $L_0$  in the 250–300 nm range at 1.5 T.

To summarize the observed trends, we can write

$$\sigma = A(\{\mathbf{r}_{\text{vortex}}\}) \exp\left(-\frac{B}{B_0(L)}\right) \times \exp\left(-\frac{T}{T_0(L)}\right), \quad (1)$$

where prefactor  $A$  represents fluctuations on the scale of a few mT due to vortex rearrangements. Importantly, the decay rate  $B_0$  is mostly independent of  $T$ , and  $T_0$  is mostly independent of  $B$  [see Fig. 2(b)], which supports our assertion that different mechanisms are responsible for the  $B$  and  $T$  dependencies.

So far we have only considered the standard deviation of  $P_{eh}$ , but it is informative to further look into its full distribution. In Fig. 4(a) we plot the histogram of  $P_{eh}$  of the 1  $\mu\text{m}$  interface over a small range of  $B = 1.4$ –1.5 T. The range  $\Delta B = 100$  mT is selected because on the one hand it covers many vortex rearrangements; on the other hand  $\Delta B = 100$  mT  $\ll B_0 = 0.85$  T so that the systematic reduction of the amplitude  $\exp(-\Delta B/B_0)$  is negligible. The measured histogram is clearly peaked around zero, which contrasts with the uniform (rectangular) distribution predicted in Ref. [23].

While the distribution gets narrower with increasing temperature, the shape remains roughly the same, as seen in Fig. 4(b), which replots the histograms with the  $x$ -axis

normalized by their widths. These widths are plotted in the right inset of Fig. 4(b) (circles), together with an exponential line whose slope is determined by the constant  $T_0$  obtained from Fig. 2(b). The decay of the distribution widths with temperature agrees with the line, except for the highest temperature, which is likely limited by the noise in determining  $P_{eh}$ .

The distribution shape is expected to be influenced by the positions of the vortices, resulting in the random spread of  $A(\{\mathbf{r}_{\text{vortex}}\})$  [see Fig. 1(c), and S6 in [29]]. To evaluate and eventually eliminate the variations of the  $P_{eh}$  due to vortex rearrangements, we normalize each  $P_{eh}(V_G)$  trace by its standard deviation calculated for the same  $B$ . The resulting histograms of  $P_{eh}(V_G)/\sigma$  are shown in Fig. S7 [29] and appear to have a roughly triangular shape similar to those in Fig. 4. We conclude that vortex rearrangements, while quite noticeable, cannot explain the shape of the  $P_{eh}$  distribution.

In principle, the nonrectangular distribution of  $P_{eh}$  may point to factors that have not yet been fully understood: spin-orbit coupling, interface transparency, contact doping, or specifics of graphene band structure. In particular, if the Andreev conversion probabilities for the two spins states are sufficiently different, the convolution of their rectangular distributions would produce a triangle. Regardless of these mechanisms, the distribution of the  $P_{eh}$  should be strongly shaped by the particle losses. In the Supplemental Material [29], we present a toy model which qualitatively reproduces the observed distribution (Fig. S8). In the model, the CAES propagation along the interface is represented by the evolution of a  $2 \times 2$  electron-hole density matrix, represented by a trajectory on the Bloch sphere [23]. The key insight of the model is that the losses have uneven effect on different trajectories: Most of the trajectories undergo heavy losses, resulting in a small  $P_{eh}$  that contributes to the center of the distribution peak. The rare trajectories which suffered less losses contribute to the tails of the distribution.

In summary, we have reported a systematic investigation of the Andreev conversion probability for the CAES states propagating along a superconducting contact in the QH regime. The main result is the near decoupling of the  $P_{eh}$  dependencies on the temperature and magnetic field, i.e., Eq. (1). The dependence on magnetic field,  $A(\{\mathbf{r}_{\text{vortex}}\}) \exp(-B/B_0)$  describes the particle loss, where the characteristic field  $B_0$  decays with the interface length  $L$  in agreement with the theoretical prediction [23]. It is natural to associate the temperature dependence,  $\exp(-T/T_0)$ , with decoherence, which is found to be much more efficient compared to the conventional edge states in quantum Hall interferometers. Indeed, the relative phase of the electron and hole components of the CAES should be strongly influenced by the fluctuations of electric environment. Understanding and controlling the loss and decoherence processes is crucial for the development of future devices that utilize superconducting correlations in chiral states.

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- [29] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/PhysRevLett.131.176604>, which includes Ref. [30] for the additional figures and notes on (i) the device fabrication and measurement methods, (ii) the comparison to the data of a nonsuperconducting interface, (iii) the data measured in additional samples, (iv) the uniform temperature dependence of  $P_{eh}(V_G)$ , (v) the variation of signal due to rearrangement of vortices, (vi) the histograms of normalized  $\sigma$  and (vii) the toy model for the distribution of  $P_{eh}$ .
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