

Aharonov-Bohm interferometer in inverted-band pn junctionsYuhao Zhao,^{1,*} Oded Zilberberg² and Antonio Štrkalj^{3,†}¹*Institute for Theoretical Physics, ETH Zurich, 8093 Zurich, Switzerland*²*Department of Physics, University of Konstanz, D-78457 Konstanz, Germany*³*Department of Physics, Faculty of Science, University of Zagreb, Bijenička c. 32, 10000 Zagreb, Croatia*

(Received 29 October 2024; accepted 7 March 2025; published 7 May 2025)

Inverted-band pn junctions in two-dimensional materials offer a promising platform for electron optics in condensed matter, as they allow to manipulate and guide electron beams without the need for spatial confinement. In this work, we propose the realization of an Aharonov-Bohm (AB) interferometer using such pn junctions. We observe AB oscillations in numerically obtained conductance and analytically identify the conditions for their appearance by analyzing the scattering processes at the pn interface. To support experimental implementation, we also consider junctions with a graded interface, where the potential varies smoothly between the p - and n -doped regions. Our results reveal an abrupt change in the AB-oscillation frequency, which we attribute to a distinct transition in the hybridization across the interface. We verify that our reported AB oscillations are robust to realistic disorder and temperature decoherence channels. Our study paves the way for realizing the Aharonov-Bohm effect in bulk mesoscopic systems without the need for external spatial confinement, offering new possibilities in electron optics.

DOI: [10.1103/PhysRevResearch.7.L022030](https://doi.org/10.1103/PhysRevResearch.7.L022030)

The Aharonov-Bohm (AB) effect is one of the cornerstone phenomena highlighting the nonlocal nature of quantum mechanics [1–3]. The effect arises when two quantum trajectories enclose a finite magnetic flux, making their interference dependent on the external magnetic field [4]. On a microscopic scale, the AB effect has been observed in a wide range of systems, and plays a pivotal role in revealing quantum phenomena in transport that strongly depends on magnetic fields, e.g., in revealing weak localization through magnetoresistance measurements [5,6].

To directly realize an AB interferometer in a mesoscopic system, a variety of two-dimensional platforms have been employed [7–14], where clear oscillations in transport as a function of an external magnetic field were observed. The primary challenges in observing such AB oscillations lie in the ability to prepare well-defined electron paths while maintaining coherence between electrons traveling along them [7,15]. Guiding electrons in such a manner is inherently difficult and has traditionally been achieved through spatial confinement using etching techniques [9,13], gating [11,16], or by employing strong magnetic fields to confine electrons along the one-dimensional edges of a quantum Hall droplet [14,17–21].

Recently [22], we proposed a protocol for directing electrons using inverted-band pn junctions, without relying on external spatial confinement of the electron trajectories. Such pn junctions filter the electronic trajectories via a Klein tunnelinglike effect, allowing passage based on the incident angle at the pn interface [22]. Unlike Klein tunneling in Dirac materials [23], the filtering occurs at finite incident angles that are readily controlled by tuning the chemical potential of the junction. Crucially to this work, a narrow window of negative refraction scattering processes—involving both particlelike and holelike Fermi surfaces—manifests and generates interfering closed paths for the electron trajectories.

In this work, we propose how to realize an Aharonov-Bohm interferometer using a two-dimensional inverted-band pn junction. The AB effect in our setup arises from two key ingredients: (i) negative refraction between distinct Fermi surfaces in the p - and n -doped regions, which creates closed loops of different electronic trajectories between pointlike source and drain leads; and (ii) angular filtering of trajectories at the pn interface, resulting in a limited number of trajectories that coherently interfere upon reaching the drain lead. In the presence of a weak perpendicular magnetic field, we predict that clear AB oscillations appear due to nonlocal interference. We numerically demonstrate the effect for a realistic experimental setting by adapting the band structure of InAs/GaAs [24]. We further show that these oscillations persist even in the situation when the pn interface has a finite length. Interestingly, we find that the frequency of the AB oscillations changes abruptly beyond a certain interface length, which we attribute to a shift in the maximal transmission angle due to a transition of the interface to an opaque barrier. Our work brings magnetic control to electron optics applications in two-dimensional mesoscopic devices.

*Contact author: yzhao.phy@gmail.com†Contact author: astrkalj@phy.hr

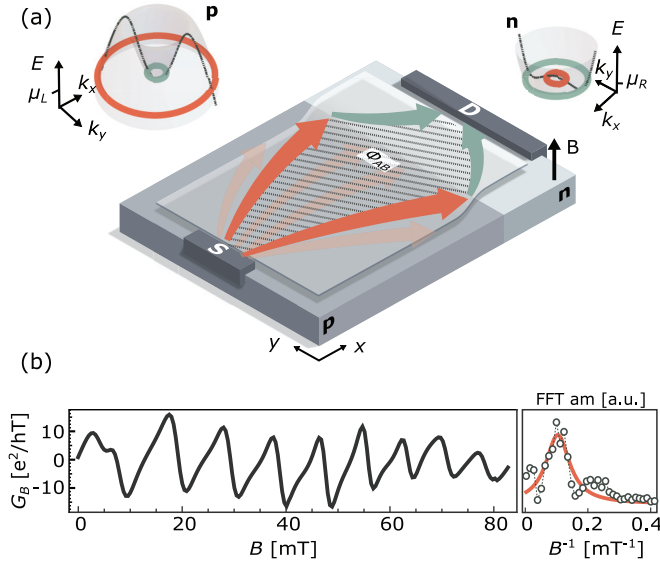


FIG. 1. Device and Aharonov-Bohm oscillations. (a) A sketch of an inverted band pn junction subject to an out-of-plane magnetic field $\mathbf{B}$. Conductance is measured between the source lead (S) and the drain lead (D). Insets: sombrero-like bandstructure for both p - and n -doped regions with marked chemical potentials $\mu_{L/R}$, cf. Eq. (1). Red (green) circles and arrows denote holelike (electronlike) branches and trajectories, respectively. (b) Left: Numerically obtained magnetodifferential conductance $G_B = dG/dB$, cf. Eq. (2). Right: Fast Fourier transform of the pattern on the left (dots) with a fitted Lorentzian (orange line). In all figures of the paper, we use the following parameters: $\mathcal{A} = 16$ meV nm, $\mathcal{M}_0 = -30$ meV, $\mathcal{M}_1 = 400$ meV nm², and $\mathcal{M}_2 = 700$ meV nm². The p - and n -doped regions are achieved by setting $\mu_L = -28$ meV and $\mu_R = 33$ meV.

Model. To model the inverted band pn junction (see Fig. 1), we consider the following 2D spinless Hamiltonian [22,24]:

$$\mathcal{H}_{L/R} = (\mathcal{M}_0 + \mathcal{M}_2 \mathbf{k}^2) \sigma_z + (\mathcal{M}_1 \mathbf{k}^2 - \mu_{L/R}) \mathbb{1}_{2 \times 2} + \mathcal{A} \sigma_x, \quad (1)$$

where σ_i are Pauli matrices acting on two bands, and $\mathbf{k}$ denotes a 2D wave vector. The bare, uncoupled energy dispersions are particle- and holelike in nature with respective effective masses $\mathcal{M}_1 \pm \mathcal{M}_2$, where we assume $\mathcal{M}_2 > \mathcal{M}_1$. Particle-hole symmetry is maintained for $\mathcal{M}_1 = 0$. Such bare energy dispersions overlap with magnitude $\mathcal{M}_0$ and hybridize with coupling strength $\mathcal{A}$. The local chemical potential $\mu_{L/R}$, experimentally tunable by backgates [24], takes different values in the p -doped and n -doped regions. We assume these regions to respectively reside on the left- and right-hand sides of the pn interface, i.e., $\mu_L < 0$ and $\mu_R > 0$. Diagonalizing the Hamiltonian (1) yields the two dispersions $E(\mathbf{k}) = \mathcal{M}_1 \mathbf{k}^2 - \mu_{L/R} \pm \sqrt{(\mathcal{M}_0 + \mathcal{M}_2 |\mathbf{k}|^2)^2 + \mathcal{A}^2}$. In the inverted-band regime, when $\mathcal{M}_0 < 0$, the resulting band structure resembles two sombreros facing one another, and separated by a gap, see Fig. 1(a). We label the spinor eigenstates of (1) as $\psi_{p/h}^{L/R, \pm}$, where L/R refers to the left- and right-hand side of the junction and p/h indicates the particle- or holelike Fermi surface, cf. Fig. 1(a). The superscript $\pm$ denotes the direction of group velocity relative to the pn interface. In the following, we consider the pn junction when the chemical

potentials, $\mu_{L/R}$, lie within the band-overlap energy interval of each sombrero. Specifically, we are interested in the response of the junction to an external perpendicular magnetic field, $\mathbf{B} = B\hat{z}$.

Sharp interface setup. To numerically probe the transport properties of the system (1), we incorporate the aforementioned junction into the numerical toolbox Kwant [25]. To simulate a realistic transport experiment, we draw the junction along the x direction, and include two metallic leads on the junction's ends. These source (S) and drain (D) leads are positions on the left and right sides, respectively, and have finite widths in the y direction, which are much shorter than the sample size, see Fig. 1(a). The spatial distance between the junction interface and the leads is chosen carefully, such that injected holelike states from S maximally refocus in D as particlelike states [26]. Similarly, the spatial extent of the junction in the y direction is chosen such that the refocusing scattering states will not hit the boundaries. To remove background noise from finite-size effects, we attach additional absorbing leads along the top and bottom boundaries. The system's dimensions we use are taken relative to realistic Fermi wavelengths of InAs/GaSb inverted-band semiconductors [24,26]. The perpendicular magnetic field $\mathbf{B}$ is introduced using standard Peierls substitution [27]. We then study the two-terminal conductance between the leads by numerically evaluating the Landauer-Buttiker formula [28]

$$G = \frac{e^2}{h} \sum_{n=1}^N \sum_{m=1}^M \text{Tr}[t_{nm}^\dagger t_{nm}], \quad (2)$$

where t_{nm} is the transmission coefficient from the n -th mode in the source lead to the m -th mode in the drain lead. The total numbers of modes N and M in the leads are determined by the widths of the leads.

Main result. In Fig. 1(b), we plot the resulting two-terminal conductance for magnetic fields ranging between 0 ~ 80 mT (cf. extended magnetic field range in the Supplemental Material [26]). We find that the magnetodifferential conductance $G_B \equiv dG/dB$ exhibits clear periodic oscillations with an average period $\Delta B \approx 8.43$ mT. As we show below, such oscillations in G_B arise due to a mesoscopic AB effect whose origin lies in negative refraction at the interface. This is the main result of our work. We emphasize that such an emergence of AB oscillations in a simple 2D pn junction that does not rely on any external spatial confinement is surprising [4].

To explain the mechanism of the oscillations in our system, we first notice that, in our regime of weak magnetic fields, the trajectories of the propagating states are only weakly affected by the magnetic field, i.e., they remain describable by “ray tracing” that emanate from the source in the p - and n -doped regions, separately [26]. Hence, in this limit, AB oscillations can emerge only when the following two prerequisites are satisfied: (i) the interface between the p - and n -doped region facilitates negative refraction that refocuses the trajectories at the drain lead where they interfere, and (ii) only a narrow distribution of incident angles participates in the negative refraction, such that the interference pattern survives decoherence via averaging over the contributions of many incident trajectories that carry different phases. Shown below, the latter

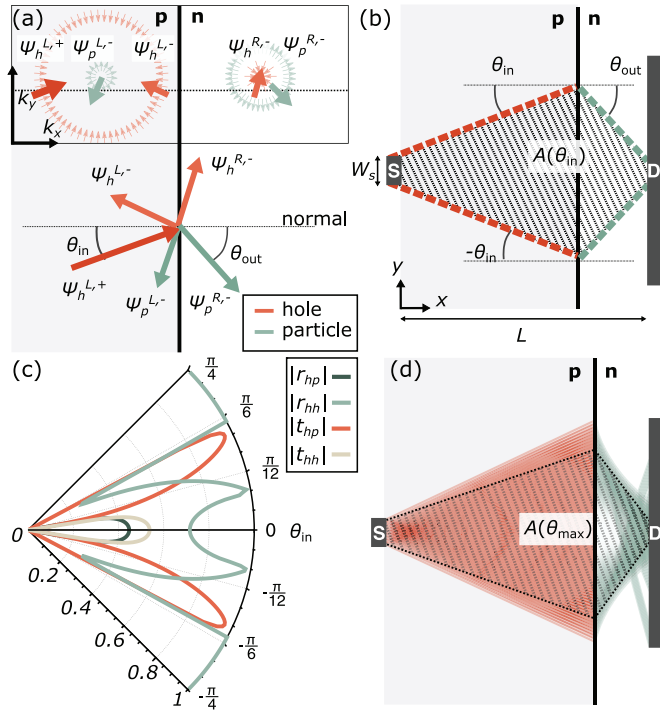


FIG. 2. Scattering processes, negative refraction, and interfering paths. (a) Upper panel: In reciprocal space, the scattering states and corresponding angles for each k_y are determined by their location on the Fermi surface. Lower panel: In real space, the scattering angles are defined with respect to the normal (dashed line) of the pn interface (solid black line). The angles θ_{in} and θ_{out} mark the incident angle of a hole-like state and the outgoing angle of a transmitted particle-like state. (b) Two trajectories of a hole-to-particle transmission scattering process with $\pm\theta_{in}$ that enclose a finite shaded area $A(\theta_{in})$, cf. Eq. (3). (c) Scattering amplitudes as a function of the incident angle for all processes involving an incident hole-like state. (d) Trajectories of the incident hole-like state injected from the source lead, with their opacity determined by the scattering amplitudes in (c). The maximal transmitted angle θ_{max} and the finite width of S define the enclosed $A(\theta_{max})$ between the source and drain lead. Here we omit trajectories of the transmitted hole-like state in the n -doped region for clarity.

will be the result of filtering by the transmission coefficients at the pn interface.

To demonstrate (i), we consider an incident hole-like spinor $\psi_h^{L,+}$ from the source with momentum $\mathbf{k} = (k_x, k_y)$, see Fig. 2(a). We assume that the pn junction is homogeneous along the y direction, such that k_y is conserved in all scattering processes. Thus, we can determine the outgoing states $\psi_{h/p}^{L/R,-}$ and their corresponding momenta $(k_{p/h}^{L/R,-}, k_y)$ from momentum conservation at the Fermi surface, cf. band structures in insets of Fig. 2(a), and Ref. [22]. Note that the incident and outgoing states are defined with respect to the pn interface, as illustrated in Fig. 2(a). The propagating direction of each scattering state is given by its group velocity $\mathbf{v} = \nabla_{\mathbf{k}} E(\mathbf{k})|_{\mathbf{k}=(k_{p/h}^{L/R,-}, k_y)}$, see arrows in upper panel of Fig. 2(a). Crucially, the hole-to-particle transmission scattering process exhibits negative refraction at the interface, i.e., for $\psi_h^{L,+}$ with a positive incident angle $\theta_{in} = \sin^{-1}(k_y/|\mathbf{k}|) > 0$, the

outgoing particlelike spinor $\psi_p^{R,-}$ in the n -doped region propagates at a negative angle $\theta_{out} < 0$. This trajectory can interfere at the drain lead with another trajectory where $\psi_h^{L,+}$ impinges on the pn interface at $-\theta_{in}$, see Fig. 2(b). Crucially, the two interfering paths encircle a kite-shaped area

$$A(\theta_{in}) = L^2 \tan(\theta_{in}) \frac{R_{pn}}{R_{pn} + 1} + A_{lead}, \quad (3)$$

where L is the distance between S and D . The value of R_{pn} represents the ratio between the length of the p - and n -doped region, which is chosen according to the focal point of trajectories [26], see also Fig. 2(b). Finally, $A_{lead} = LW_s$ accounts for the added area introduced by the finite width W_s of the source lead.

To check if (ii) is satisfied, we employ a scattering matrix formalism [26] to calculate the transmission t_{hp} and t_{hh} and reflection r_{hp} and r_{hh} scattering coefficients as a function of the incident angle, θ_{in} , where we assume an incident holelike spinor (first subscript = h) and consider all four possible outgoing cases, see Fig. 2(c). At most angles the dominant outgoing process is a reflection of a hole, r_{hh} , making the junction very opaque. Interestingly, we find that t_{hp} , which leads to negative refraction, maximizes in a narrow range around a finite incident angle $|\theta_{in}| \approx \theta_{max} = 5\pi/36$ [26]. In other words, electrons transmitted through the pn junction are mostly carried by trajectories incident around θ_{max} , while the other angles are screened by the pn interface. We demonstrate this effect by drawing the ray trajectories of the holelike states injected from the source and transmitted as particlelike states into the drain lead, see Fig. 2(d). The drawn trajectories' opacity is taken according to their transmission amplitude at the interface, cf. Fig. 2(c). Thus, we find a large depleted area due to the transmission filtering at small angles and high density only in a narrow scattering angle region. Note that by taking only the two trajectories with the maximal t_{hp} into account, defined by $\theta = \pm\theta_{max}$ and marked by dashed lines in Fig. 2(d), we can use Eq. (3) to obtain $A(\theta_{max}) = 4.69 \times 10^5 \text{ nm}^2$ for $W_s = 80 \text{ nm}$ and $R_{pn} = 131/44$ in our setup.

Each scattering trajectory follows a path $\mathcal{P}$ and accumulates a dynamical phase $\phi_D = \exp(\int_{\mathbf{x} \in \mathcal{P}} d\mathbf{l} \cdot \mathbf{k})$. To these paths, we can introduce the perpendicular magnetic field as a vector potential $\mathbf{A} = (-By/2, Bx/2, 0)$, where it may impact the scattering amplitudes in three ways: (i) bend the trajectories due to a Lorentz force, (ii) alter the scattering angles and amplitudes at the junction, and (iii) introduce an Aharonov-Bohm (AB) phase to the interfering paths via minimal coupling, $\exp(\int_{\mathbf{x} \in \mathcal{P}} d\mathbf{l} \cdot \mathbf{A})$. Neglecting (i) and (ii) due to the weak magnetic fields we used, we can divide the sum in Eq. (2) into terms that do not include interfering paths (where the accumulated phase is irrelevant) and interfering scattering trajectories that are impacted by the magnetic flux through the area A . We can, moreover, assume that the dominant magnetic-field dependent scattering process involves the paths at $\theta_{in} = \pm\theta_{max}$ [26], leading to

$$G_B \approx \frac{e^2}{h} |t_{hp}(\theta_{max})|^2 \frac{d}{dB} [1 + \cos(\phi_{tot})], \quad (4)$$

where $\phi_{tot} = eBA(\theta_{max})/h$ is the flux threading the enclosed area $A(\theta_{max})$. The magnetic response (4), therefore, exhibits a

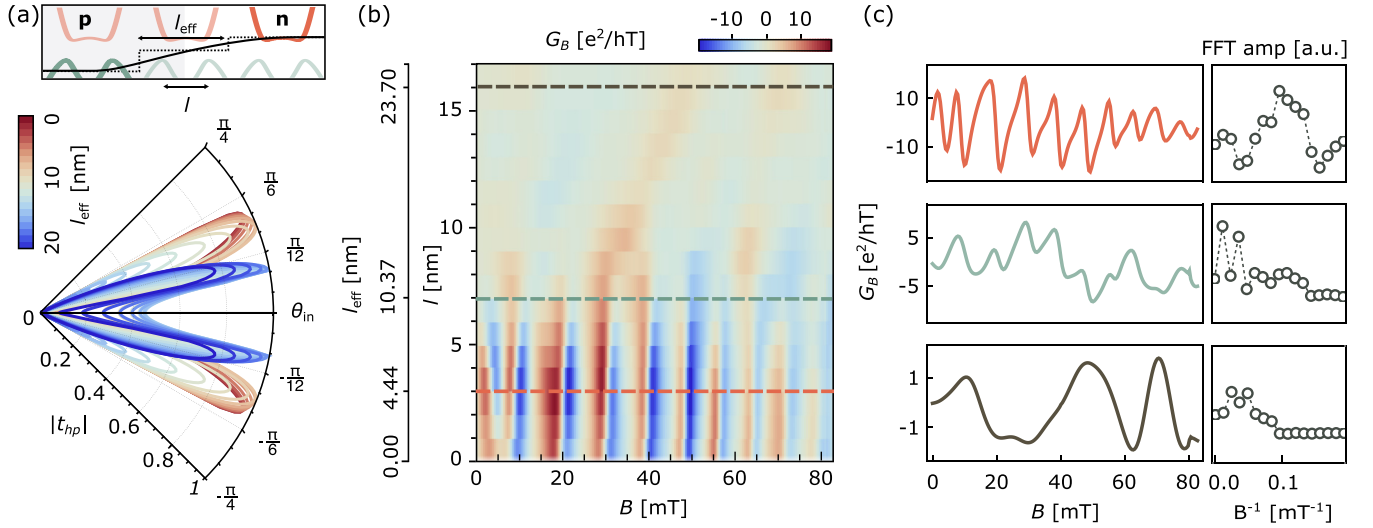


FIG. 3. Smooth interface. (a) Upper panel: A spatially varied chemical potential (solid line) induces an effective barrier defined by the length over which the Fermi level is inside the gap, leading to an effective two-interface junction (dotted line). Lower panel: Transmission amplitude of a hole-to-particle $|t_{hp}|$ scattering as a function of the incident angle θ_{in} and the effective barrier length l_{eff} . (b) Numerically obtained G_B as a function of the magnetic field B and the (effective) barrier length $(l_{\text{eff}}) l$. (c) Left: Cuts of (b) at $l = 3, 7, 16$ nm (red, green, black), cf. dashed lines. Right: Fast Fourier transform of the patterns.

period $\Delta B = h[eA(\theta_{\text{max}})]^{-1}$ that is solely determined by the AB phase. In our case, we obtain a period of $\Delta B = 8.82$ mT, which is in good agreement with the one extracted from the numerics, cf. Fig. 1(b).

Graded pn interface. Perfectly sharp pn interfaces cannot be realized in experiments. We, therefore, proceed to check how the predicted AB oscillations behave in the presence of a graded interface. We model it by introducing a spatially modulated chemical potential $\mu(x) = \mu_R + (\mu_L - \mu_R)[1 - \tanh(x/l)]/2$, where l is its characteristic length, see in Fig. 3(a). Importantly, this interpolation introduces an effective gapped region between the p and n regions. We numerically calculate G_B as a function of l and show it in Fig. 3(b). We observe a clear change in the oscillating pattern at $l \approx 7$ nm. More quantitatively, the frequency of oscillation in G_B reduces for smoother interfaces, cf. three cuts at $l = 3$ nm, 7 nm, and 16 nm in Fig. 3(c).

To gain a deeper understanding of the observations discussed above, we analytically describe an effective potential barrier induced by the smooth interface by employing an effective two-interface model. As the Fermi level leaves the upper/lower band at energies $\Lambda_{\pm} = (-\mathcal{M}_1\mathcal{M}_0 \pm A\sqrt{\mathcal{M}_2^2 - \mathcal{M}_1^2})/\mathcal{M}_2$, the barrier between the p - and n -doped regions has an effective length $l_{\text{eff}} = x_+ - x_-$ with $x_{\pm} = \text{arctanh}(1 - 2(\Lambda_{\pm} - \mu_R)/(\mu_L - \mu_R))/l$. In this simplified model, we assume the interfaces between the three regions to be sharp; see the dashed line in the upper panel of Fig. 3(a). Using the S-matrix formalism, we once more calculate the transmission amplitude t_{hp} as a function of the incident angle and effective barrier length l_{eff} [26] and show the result in the lower panel of Fig. 3(a). For a short barrier $l_{\text{eff}} < 10$ nm, we obtain a similar behavior as in the previous case of a vanishing barrier, with negative refraction centered around $\theta_{\text{max}} \approx 5\pi/36$. Interestingly, with increasing l_{eff} , we find a crossover in the behavior of $|t_{hp}|$, i.e., the

maximal transmitted angle θ_{max} undergoes an abrupt transition at $l_{\text{eff}} \approx 10$ nm. As a result, for a longer interface $l_{\text{eff}} > 10$ nm, $|t_{hp}|$ maximizes at a smaller angle of $\theta'_{\text{max}} \approx \pi/18$, which implies a smaller enclosed area $A(\theta'_{\text{max}})$, and a different AB-oscillation period that follows from Eq. (4), which is in agreement with the numerical results, cf. Fig. 3. We attribute this transition to a competition between the length of the barrier and the decay length of evanescent waves penetrating the barrier from the interfaces [29]: the system enters the tunneling limit when the electrons only tunnel through the overlapping evanescent tails, resulting in an abrupt change of the θ_{max} [26]. Moreover, we find that the oscillating pattern persists to larger magnetic fields for small θ_{max} , which is favorable to experimental realizations [26].

In conclusion, we propose the realization of an AB interferometer using negative refraction in a two-dimensional inverted-band pn junction. A key difference, compared with common AB interferometers, is that we do not impose external spatial confinement on the trajectories of electrons; instead, the confinement emerges naturally due to the band structure's characteristic sombrero shape. Specifically, using the S-matrix formalism, we find that the negative refraction transmission amplitudes are facilitated by hole-to-particle scattering t_{hp} , whose distribution maximizes in a narrow range of incident angles. The latter realizes an effective spatial confinement that guides electrons [22]. This emergent spatial confinement entails that all transmitted electronic trajectories exhibit a kite-shaped enclosed surface which, under the external magnetic field, gives rise to unique AB oscillations in conductance. Commonly, one would expect a reduction in the transmission due to a smooth interface [23], while, on the contrary, we find that the oscillations persist with a distinct change in the AB-oscillation period. Moreover, we tested the robustness of the AB oscillations against finite temperature and disorder [26]. We observe clear oscillations

up to temperatures of 1 K and disorder strength of several meV's, which gives strong indications that the proposed setup can be realized in realistic experimental situations. The proposed mechanism of generating AB interference is related to the sombrero hat shape of the band structure, which can be found in a plethora of systems, including s -wave and chiral p -type superconductors [30], heterostructures with strong Rashba spin-orbit interaction like HgTe/CdTe [31] and InAs/GaSb [24], two-dimensional monochalcogenides like GaSe [32], and transition metal dichalcogenides like WTe₂ [33]. Moreover, the Bernal bilayer graphene, which is known for its tunability of the sombrero-hat bandstructure under a

biased voltage [34–36], suggests potential realizations of AB interferometer with graphene-based material. With all of the aforementioned, we think that the physics discussed in this work can be experimentally tested in current state-of-the-art mesoscopic platforms.

This work was supported by the Deutsche Forschungsgemeinschaft (DFG) through Project No. 449653034, and ETH research Grant ETH-28 23-1. The work of A.Š. is supported by the European Union's Horizon Europe research and innovation programme under the Marie Skłodowska-Curie Actions Grant Agreement No. 101104378.

- [1] W. Ehrenberg and R. E. Siday, The refractive index in electron optics and the principles of dynamics, *Proc. Phys. Soc. B* **62**, 8 (1949).
- [2] Y. Aharonov and D. Bohm, Significance of electromagnetic potentials in the quantum theory, *Phys. Rev.* **115**, 485 (1959).
- [3] Y. Aharonov and D. Bohm, Further considerations on electromagnetic potentials in the quantum theory, *Phys. Rev.* **123**, 1511 (1961).
- [4] T. Ihn, *Semiconductor Nanostructures: Quantum States and Electronic Transport* (Oxford University Press, Oxford, UK, 2009).
- [5] E. Akkermans and G. Montambaux, *Mesoscopic Physics of Electrons and Photons* (Cambridge University Press, Cambridge, UK, 2007).
- [6] W. Chen, H.-Z. Lu, and O. Zilberberg, Weak localization and antilocalization in nodal-line semimetals: Dimensionality and topological effects, *Phys. Rev. Lett.* **122**, 196603 (2019).
- [7] A. van Oudenaarden, M. H. Devoret, Y. V. Nazarov, and J. E. Mooij, Magneto-electric Aharonov–Bohm effect in metal rings, *Nature (London)* **391**, 768 (1998).
- [8] A. Bachtold, C. Strunk, J.-P. Salvetat, J.-M. Bonard, L. Forró, T. Nussbaumer, and C. Schönenberger, Aharonov–Bohm oscillations in carbon nanotubes, *Nature (London)* **397**, 673 (1999).
- [9] S. Russo, J. B. Oostinga, D. Wehenkel, H. B. Heersche, S. S. Sobhani, L. M. K. Vandersypen, and A. F. Morpurgo, Observation of Aharonov–Bohm conductance oscillations in a graphene ring, *Phys. Rev. B* **77**, 085413 (2008).
- [10] M. Kremer, I. Petrides, E. Meyer, M. Heinrich, O. Zilberberg, and A. Szameit, A square-root topological insulator with non-quantized indices realized with photonic Aharonov–Bohm cages, *Nat. Commun.* **11**, 907 (2020).
- [11] S. Iwakiri, A. Mestre-Torà, E. Portolés, M. Visscher, M. Perego, G. Zheng, T. Taniguchi, K. Watanabe, M. Sigrist, T. Ihn, and K. Ensslin, Tunable quantum interferometer for correlated moiré electrons, *Nat. Commun.* **15**, 390 (2024).
- [12] Y. Kawaguchi, D. Smirnova, F. Komissarenko, S. Kiriushechkina, A. Vakulenko, M. Li, A. Alù, and A. B. Khanikaev, Pseudo-spin switches and Aharonov–Bohm effect for topological boundary modes, *Sci. Adv.* **10**, eadn6095 (2024).
- [13] J. Kim, H. Dev, R. Kumar, A. Ilin, A. Haug, V. Bhardwaj, C. Hong, K. Watanabe, T. Taniguchi, A. Stern, and Y. Ronen, Aharonov–Bohm interference and statistical phase-jump evolution in fractional quantum Hall states in bilayer graphene, *Nat. Nanotechnol.* **19**, 1619 (2024).
- [14] S. Biswas, H. K. Kundu, R. Bhattacharyya, V. Umansky, and M. Heiblum, Anomalous Aharonov–Bohm interference in the presence of edge reconstruction, *Phys. Rev. Lett.* **132**, 076301 (2024).
- [15] O. Zilberberg, A. Romito, and Y. Gefen, Charge sensing amplification via weak values measurement, *Phys. Rev. Lett.* **106**, 080405 (2011).
- [16] S. Iwakiri, F. K. De Vries, E. Portolés, G. Zheng, T. Taniguchi, K. Watanabe, T. Ihn, and K. Ensslin, Gate-defined electron interferometer in bilayer graphene, *Nano Lett.* **22**, 6292 (2022).
- [17] K. v. Klitzing, G. Dorda, and M. Pepper, New method for high-accuracy determination of the fine-structure constant based on quantized Hall resistance, *Phys. Rev. Lett.* **45**, 494 (1980).
- [18] B. Özyilmaz, P. Jarillo-Herrero, D. Efetov, D. A. Abanin, L. S. Levitov, and P. Kim, Electronic transport and quantum Hall effect in bipolar graphene p - n - p junctions, *Phys. Rev. Lett.* **99**, 166804 (2007).
- [19] F. Amet, J. R. Williams, K. Watanabe, T. Taniguchi, and D. Goldhaber-Gordon, Selective equilibration of spin-polarized quantum Hall edge states in graphene, *Phys. Rev. Lett.* **112**, 196601 (2014).
- [20] F. Nichele, H. J. Suominen, M. Kjaergaard, C. M. Marcus, E. Sajadi, J. A. Folk, F. Qu, A. J. A. Beukman, F. K. de Vries, J. v. Veen, S. Nadj-Perge, L. P. Kouwenhoven, B.-M. Nguyen, A. A. Kiselev, W. Yi, M. Sokolich, M. J. Manfra, E. M. Spanton, and K. A. Moler, Edge transport in the trivial phase of InAs/GaSb, *New J. Phys.* **18**, 083005 (2016).
- [21] B.-M. Nguyen, A. A. Kiselev, R. Noah, W. Yi, F. Qu, A. J. A. Beukman, F. K. de Vries, J. van Veen, S. Nadj-Perge, L. P. Kouwenhoven, M. Kjaergaard, H. J. Suominen, F. Nichele, C. M. Marcus, M. J. Manfra, and M. Sokolich, Decoupling edge versus bulk conductance in the trivial regime of an InAs/GaSb double quantum well using Corbino ring geometry, *Phys. Rev. Lett.* **117**, 077701 (2016).
- [22] Y. Zhao, A. Leuch, O. Zilberberg, and A. Štrkalj, Electron optics using negative refraction in two-dimensional inverted-band pn junctions, *Phys. Rev. B* **108**, 195301 (2023).
- [23] V. V. Cheianov and V. I. Fal'ko, Selective transmission of Dirac electrons and ballistic magnetoresistance of n - p junctions in graphene, *Phys. Rev. B* **74**, 041403(R) (2006).
- [24] M. Karalic, A. Štrkalj, M. Masseroni, W. Chen, C. Mittag, T. Tschirky, W. Wegscheider, T. Ihn, K. Ensslin, and O. Zilberberg, Electron-hole interference in an inverted-band semiconductor bilayer, *Phys. Rev. X* **10**, 031007 (2020).

- [25] C. W. Groth, M. Wimmer, A. R. Akhmerov, and X. Waintal, Kwant: a software package for quantum transport, *New J. Phys.* **16**, 063065 (2014).
- [26] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/PhysRevResearch.7.L022030> for the details on (1) setup of the numeric simulations, (2) results for finite temperatures and the disorder, (3) analytic calculation using the scattering matrix formalism, (4) determining the effective path, the focal point and the carried phases, (5) results at the higher magnetic field, (6) comparison to an AB interferometer with the spatial confinement, (7) the fate of the injected particle-like spinor, (8) determining the angle of the maximal transmission, (9) scattering at the graded interface.
- [27] R. Peierls, Zur theorie des diamagnetismus von leitungselektronen, *Zeitschrift für Physik* **80**, 763 (1933).
- [28] M. Büttiker, Absence of backscattering in the quantum Hall effect in multiprobe conductors, *Phys. Rev. B* **38**, 9375 (1988).
- [29] J. Yao, B. Zhou, X. Zhou, X. Xiao, and G. Zhou, Anomalous quantum scattering and transport of electrons with Mexican-hat dispersion induced by electrical potential, [arXiv:2403.09319](https://arxiv.org/abs/2403.09319).
- [30] A. Štrkalj, X.-R. Chen, W. Chen, D. Y. Xing, and O. Zilberberg, Tomash oscillations as above-gap signature of topological superconductivity, *Phys. Rev. Lett.* **132**, 066301 (2024).
- [31] G. M. Minkov, A. V. Germanenko, O. E. Rut, A. A. Sherstobitov, S. A. Dvoretzki, and N. N. Mikhailov, Two-dimensional semimetal in a wide HgTe quantum well: Magnetotransport and energy spectrum, *Phys. Rev. B* **88**, 155306 (2013).
- [32] X. Li, M.-W. Lin, A. A. Purotzky, J. C. Idrobo, C. Ma, M. Chi, M. Yoon, C. M. Rouleau, I. I. Kravchenko, D. B. Geohegan, and K. Xiao, Controlled vapor phase growth of single crystalline, two-dimensional gas crystals with high photoresponse, *Sci. Rep.* **4**, 5497 (2014).
- [33] X. Qian, J. Liu, L. Fu, and J. Li, Quantum spin Hall effect in two-dimensional transition metal dichalcogenides, *Science* **346**, 1344 (2014).
- [34] H. Min, B. Sahu, S. K. Banerjee, and A. H. MacDonald, *Ab initio* theory of gate induced gaps in graphene bilayers, *Phys. Rev. B* **75**, 155115 (2007).
- [35] E. V. Castro, K. S. Novoselov, S. V. Morozov, N. M. R. Peres, J. M. B. Lopes Dos Santos, J. Nilsson, F. Guinea, A. K. Geim, and A. H. Castro Neto, Biased bilayer graphene: semiconductor with a gap tunable by the electric field effect, *Phys. Rev. Lett.* **99**, 216802 (2007).
- [36] E. McCann, D. S. Abergel, and V. I. Fal'ko, The low energy electronic band structure of bilayer graphene, *Eur. Phys. J.: Spec. Top.* **148**, 91 (2007).