

Anomalous Fano factor as a signature of Bogoliubov Fermi surfaces

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Noise spectroscopy is a key technique to investigate the nature and dynamics of charge carriers in superconductors. The recently discovered superconducting hybrids with Bogoliubov Fermi surfaces exhibit a particularly intriguing and rich charge dynamics, as their charge carriers consist of both Cooper pairs and an extensive number of Bogoliubov quasiparticles. Motivated by this, we compute the noise spectra of Bogoliubov Fermi surfaces and identify their key signatures in the differential conductance and the Fano factor. Specifically, we consider a semiconductor/superconductor hybrid device with an in-plane magnetic field, which exhibits several Bogoliubov Fermi surfaces. The number and orientation of the Bogoliubov Fermi surfaces in this device can be readily controlled by the applied magnetic field, which in turn alters the noise signal. In particular, we find that the Fano factor exhibits a reduced value, substantially lower than two, whenever the charge dynamics is governed by a large number of Bogoliubov quasiparticles. Using experimentally relevant parameters, we make a number of specific predictions for the noise spectra, that can be used as direct evidence of Bogoliubov Fermi surfaces. In particular, we find that the Fano factor as a function of magnetic field and spin-orbit coupling exhibits characteristic discontinuities at the transition lines that separate phases with different numbers of Bogoliubov Fermi surfaces.

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Introduction. Since the discovery of unconventional superconductors in the 1980s [1–3], the study of low-energy Bogoliubov quasiparticles (BQPs) has been an important research topic with both fundamental and applied interest [4]. Usually, unconventional superconductors exhibit only point nodes or line nodes in their pairing gap, with dimension strictly smaller than the normal-state Fermi surface, due to the loss of condensation energy. However, recent findings have shown that there can also exist two-dimensional nodes in the superconducting gap, which form a Fermi surface of BQPs whose dimension is equal to the normal-state Fermi surface. For instance, Bogoliubov Fermi surfaces (BFSs) can exist in multiband superconductors with broken time-reversal symmetry [5–10], in nodal superconductors with external Zeeman potentials [11–13], and in superconductors carrying large supercurrents [14,15]. In addition, BFSs can be created in superconducting hybrids, such as semiconductor-superconductor heterostructures with strong spin-orbit coupling (SOC) and an external Zeeman field [16,17], or at the interface between an *s*-wave superconductor

and a halfmetal [18–20]. The recent experimental observation of BFSs in superconductor hybrids [15,16] has further stimulated this field.

The presence of BFSs with their extensive number of BQPs leads to a large density of states at zero energy, which can be detected by a number of experimental probes. For example, the large zero-energy density of states can be directly observed in tunneling conductance [13,21] and specific heat [13,21], and moreover leads to characteristic signatures in quasiparticle interference patterns [15]. While these experimental probes detect the presence of zero-energy quasiparticles, they do not provide information about their nature and dynamics. This is where noise spectroscopy comes into play, which can measure the fundamental charge of the quasiparticles as well as their dynamics. For example, the current-noise ratio (i.e., Fano factor) in the tunneling limit of a fully gapped superconductor is twice as large as in the normal state [22,23], indicating that the charge carriers are Cooper pairs composed of two electrons. However, in systems with BFSs, where the charge carriers consist of both Cooper pairs and BQPs, the Fano factor is expected to be smaller than two since the effective charge of the BQPs is substantially smaller than that of Cooper pairs. Hence, noise spectroscopy can give valuable information about the presence of BFSs and can tell us to which degree BQPs contribute to charge transport.

Motivated by these considerations, we investigate in this Letter the Fano factor of a semiconductor/superconductor hybrid device which exhibits BFSs that can be controlled by an external magnetic field. After defining the model Hamiltonian of the considered device (Sec. II), we determine in Sec. III

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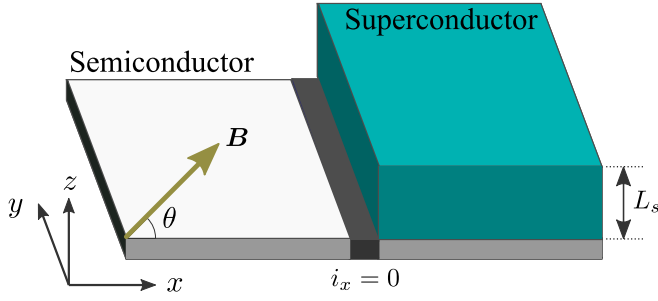


FIG. 1. Schematics of the considered semiconductor/superconductor hybrid device that exhibits Bogoliubov Fermi surfaces. An external field $\mathbf{B}$ is applied inplane and can be rotated by the angle θ . A potential barrier at $i_x = 0$ separates the superconducting from the semiconducting region.

the regions in parameter space of SOC and field strength, where BFSs exist. Interestingly, we find that with increasing field strength our device can be tuned through two phase transitions, where the number of BFSs increases from zero to two to four (Fig. 2). In Sec. IV we present our results on the differential conduction and noise spectroscopy. We observe that both the differential conductance and the Fano factor exhibit characteristic discontinuities at the phase transition lines. In the tunneling limit the value of the Fano factor reduces from two to about one, indicating that a substantial part of the charge is carried by BQPs rather than Cooper pairs. These results provide key fingerprints for the experimental identification of BFSs in the considered device.

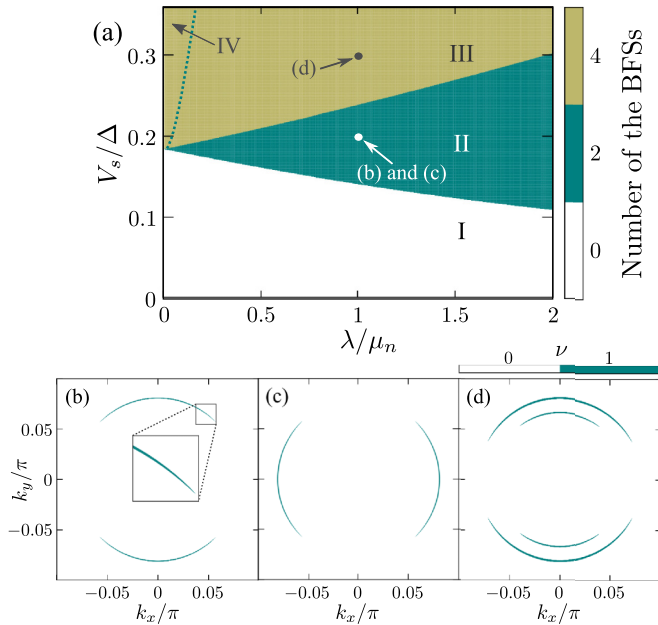


FIG. 2. (a) Phase diagram of the device shown in Fig. 1 as a function of Rashba SOC strength λ and Zeeman field V_s . Phase II has two BFSs, while phases III and IV have four BFSs. Panels (b), (c), and (d) show the determinant index ν , Eq. (5), in the two-dimensional Brillouin zone for $(\lambda/\mu_n, V_s/\Delta, \theta/\pi) = (1.0, 0.2, 0)$, $(1.0, 0.2, 0.5\pi)$, and $(1.0, 0.3, 0)$, respectively, as indicated in panel (a). The colored regions are bounded by BFSs.

Model definition. We consider a thin-film semiconductor/superconductor hybrid device with an external magnetic field applied inplane, as illustrated in Fig. 1. The superconductor is assumed to be of conventional nature with spin-singlet s -wave pairing, while the semiconductor is noncentrosymmetric with a substantial Rashba SOC, as realized in, e.g., Al/InAs heterostructures [16,24]. The semiconductor is separated into two regions by a potential barrier at $i_x = 0$, whose transparency can be controlled experimentally by applying gate voltages. Due to proximity to the parent superconductor, the semiconductor at $i_x > 0$ acquires an induced s -wave pair potential [16,24–27]. As a result, the hybrid device realizes an effective two-dimensional NS junction with tunable potential barrier and in-plane field. As we will see, the high tunability of the device allows us to study the conductance and noise signals for different BFSs configurations and junction transparencies.

Mathematically, we describe the considered device by a Bogoliubov-de Gennes (BdG) Hamiltonian on a three-dimensional cubic lattice with lattice sites $\mathbf{r}$ labeled by $\mathbf{r} = i_x\mathbf{x} + i_y\mathbf{y} + i_z\mathbf{z}$, where the three lattice vectors $\mathbf{x}$, $\mathbf{y}$, and $\mathbf{z}$ are assumed to have unit length $a_0 = 1$. We can separate the BdG Hamiltonian into three parts, $H = H_n + H_s + H_i$, representing the semiconductor, the superconductor, and the coupling between them, respectively. The Hamiltonian of the two-dimensional semiconductor at $i_z = 0$ is given by

$$H_n = \sum_{\bar{\mathbf{r}}} \sum_{\bar{\mathbf{w}}=\pm\mathbf{x},\pm\mathbf{y}} C_{\bar{\mathbf{r}}+\bar{\mathbf{w}}}^\dagger \left\{ -t_n \hat{\sigma}_0 - \frac{i\lambda}{2} \mathbf{z} \cdot (\hat{\boldsymbol{\sigma}} \times \bar{\mathbf{w}}) \right\} C_{\bar{\mathbf{r}}} + \sum_{\bar{\mathbf{r}}} C_{\bar{\mathbf{r}}}^\dagger \{ (\epsilon_n + X_b \delta_{i_x,0}) \hat{\sigma}_0 + \mathbf{V}_n \cdot \hat{\boldsymbol{\sigma}} \} C_{\bar{\mathbf{r}}}, \quad (1)$$

where $\hat{\boldsymbol{\sigma}} = (\hat{\sigma}_x, \hat{\sigma}_y, \hat{\sigma}_z)$ are the three Pauli matrices in spin space, and $\hat{\sigma}_0$ denotes the (2×2) unit matrix. Here, $C_{\bar{\mathbf{r}}} = [c_{\bar{\mathbf{r}},\uparrow}, c_{\bar{\mathbf{r}},\downarrow}]^T$ represents a spinor composed of operators $c_{\bar{\mathbf{r}},\sigma}$ ($c_{\bar{\mathbf{r}},\sigma}^\dagger$) which annihilate (create) electrons with spin σ at sites $\bar{\mathbf{r}} = i_x\mathbf{x} + i_y\mathbf{y}$ within the semiconductor. t_n denotes the nearest-neighbor hopping amplitude and $\epsilon_n = 4t_n - \mu_n$ is the onsite energy, with μ_n the chemical potential. The strength of the Rashba SOC and the potential barrier at $i_x = 0$ are denoted by λ and X_b , respectively. The in-plane magnetic field $\mathbf{B} = B(\cos \theta, \sin \theta, 0)$ induces a Zeeman potential $\mathbf{V}_n = V_n(\cos \theta, \sin \theta, 0)$ of strength $V_n = g_n \mu_B B$, where μ_B and g_n denote the Bohr magneton and the g factor of the semiconductor, respectively.

The Hamiltonian of the three-dimensional superconductor, located at $i_x > 0$ and $0 < i_z \leq L_s$, reads

$$H_s = -t_s \sum_{i_x>0} \sum_{i_y} \sum_{i_z=1}^{L_s} \sum_{\bar{\mathbf{w}}=\mathbf{x},\mathbf{y}} [D_{\bar{\mathbf{r}}+\bar{\mathbf{w}},i_z}^\dagger D_{\bar{\mathbf{r}},i_z} + \text{H.c.}] - t_s \sum_{i_x>0} \sum_{i_y} \sum_{i_z=1}^{L_s-1} (D_{\bar{\mathbf{r}},i_z+1}^\dagger D_{\bar{\mathbf{r}},i_z} + \text{H.c.}) + \sum_{i_x>0} \sum_{i_y} \sum_{i_z=1}^{L_s} D_{\bar{\mathbf{r}},i_z}^\dagger \{ \epsilon_s \hat{\sigma}_0 + \mathbf{V}_s \cdot \hat{\boldsymbol{\sigma}} \} D_{\bar{\mathbf{r}},i_z} + \sum_{i_x>0} \sum_{i_y} \sum_{i_z=1}^{L_s} \Delta (d_{\bar{\mathbf{r}},i_z,\uparrow}^\dagger d_{\bar{\mathbf{r}},i_z,\downarrow}^\dagger + \text{H.c.}), \quad (2)$$

where $D_{\bar{r},i_z} = [d_{\bar{r},i_z,\uparrow}, d_{\bar{r},i_z,\downarrow}]^T$ is a spinor composed of operators $d_{\bar{r},i_z,\sigma}$ ($d_{\bar{r},i_z,\sigma}^\dagger$) which annihilate (create) electrons with spin σ at sites $\mathbf{r} = \bar{\mathbf{r}} + i_z \mathbf{z}$ within the superconductor. t_s denotes the hopping amplitude in the superconductor, Δ represents the spin-singlet s -wave pair potential, and $\epsilon_s = 6t_s - \mu_s$ is the onsite energy, with μ_s the chemical potential. The Zeeman potential in the superconductor is given by $\mathbf{V}_s = V_s(\cos\theta, \sin\theta, 0)$, with strength $V_s = g_s \mu_B B$ and g -factor g_s . We note that we have assumed for simplicity a uniform Zeeman potential within the superconductor, which is well justified as long as the superconducting film is thinner than the London penetration depth, which can be experimentally achieved using, for example, Al films [28]. Finally, the coupling between the semiconductor and the superconductor is given by

$$H_i = -t_i \sum_{i_x > 0} \sum_{i_y} (D_{\bar{r},i_z=1}^\dagger C_{\bar{r}} + \text{H.c.}), \quad (3)$$

with amplitude t_i . This coupling induces a finite superconducting pair potential in the semiconductor segment $i_x > 0$, which we compute using lattice Greens function techniques [29,30].

For the numerical calculations, we set the parameters to $t_n = 10t_s$, $\mu_n = 0.5t_s$, $\mu_s = 3t_s$, $\Delta = 0.01t_s$, $t_i = 0.2t_s$, and $L_s = 20$. In addition, we fix the ratio of the g factors to $g_n/g_s = 5$ and express the strength of the potential barrier by the normalized value $Z = X_b/(2t_n \sin k_F)$, with $k_F = \arccos[1 - (\mu_n/2t_n)]$ the Fermi momentum of the semiconductor. We note that the choice of Δ and g_n/g_s corresponds to the experimentally relevant regime, as observed, e.g., in Al/InAs heterostructures [16,24]. In the following, we calculate the differential conductance and the Fano factor as a function of junction transparency Z , SOC strength λ , and Zeeman field V_s , for different field orientations θ . We note that with this choice V_n is given by $V_n = 5V_s$. We emphasize that our main findings do not depend on the particular parameter choices. In particular, the reduction of the Fano factor below one and the sudden drop of the Fano factor across the transition [Fig. 4] are present in any semiconductor/superconductor hybrid device with BFSs.

Phase diagram. Before calculating the conductance and noise spectra, we first determine the parameter regions where BFSs emerge in the considered device. For this purpose, we remove the semiconductor segment $i_x \leq 0$ and apply periodic boundary conditions in both x and y directions. By Fourier transforming the spinors $C_{\bar{r}}$ and $D_{\bar{r},i_z}$, as $C_{\bar{r}} = \sum_{\bar{k}} e^{i\bar{k}\cdot\bar{r}} C_{\bar{k}}$ and $D_{\bar{r},i_z} = \sum_{\bar{k}} e^{i\bar{k}\cdot\bar{r}} D_{\bar{k},i_z}$, the BdG Hamiltonian is simplified to

$$H = \frac{1}{2} \sum_{\bar{k}} \Psi_{\bar{k}}^\dagger H(\bar{k}) \Psi_{\bar{k}}, \quad (4)$$

where $H(\bar{k})$ is given in the Supplemental Material (SM) [31] and $\Psi_{\bar{k}} = [C_{\bar{k}}, D_{\bar{k},1}, \dots, D_{\bar{k},L_s}]^T$, with $C_{\bar{k}} = [C_{\bar{k}}, C_{-\bar{k}}^\dagger]^T$ and $D_{\bar{k},i_z} = [D_{\bar{k},i_z}, D_{-\bar{k},i_z}^\dagger]^T$. In the presence of time-reversal symmetry (TRS) and inversion, the emergence of BFSs can be detected by use of a $\mathbb{Z}_2$ Pfaffian index, as discussed in Refs. [6–8,32]. In the case of Hamiltonian (4), however, the Pfaffian index is no longer well defined, since TRS and inversion are broken by the magnetic field and Rashba SOC, respectively. Instead, we introduce for a given wave vector

$\bar{k} = (k_x, k_y)$ of the two-dimensional Brillouin zone (BZ) a $\mathbb{Z}_2$ determinant index ν , which can be used even in the absence of TRS and inversion. It is defined by

$$(-1)^\nu = \text{sgn}[\det H(\bar{k}) \times \det H(\bar{k}_0)], \quad (5)$$

where $\bar{k}_0$ is the (arbitrarily chosen) origin of the BZ. We note that ν is defined modulo two, with $\nu = 1$ indicating that the energy gap of the BdG Hamiltonian $H(\bar{k})$ closes at a momentum between $\bar{k}$ and the BZ origin $\bar{k}_0$. Hence, two-dimensional regions with $\nu = 1$ in the BZ are bounded by a gapless line that forms a BFS. This is demonstrated in Figs. 2(b)–2(d) where we plot ν in the two-dimensional BZ for three different parameter choices indicated in Fig. 2(a), namely $(\lambda/\mu_n, V_s/\Delta, \theta/\pi) = (1.0, 0.2, 0)$, $(1.0, 0.2, 0.5\pi)$, and $(1.0, 0.3, 0)$, respectively. We observe pairs of thin banana-shaped regions where $\nu = 1$, which are bounded by BFSs.

The orientation of these BFSs is controlled by the field angle θ . That is, with increasing θ the BFSs rotate counterclockwise from the 12 o'clock and 6 o'clock positions for $\theta = 0$ [Fig. 2(b)] to the 3 o'clock and 9 o'clock positions for $\theta = \pi/2$ [Fig. 2(c)]. The length of the banana-shaped BFSs can be tuned by the Zeeman field V_s . That is, with increasing V_s the BFSs become longer and longer. For sufficiently large V_s a phase transition point is reached, at which two more point-like BFSs appear, which are inflated into banana shapes, by further increasing V_s . As a function of Zeeman field V_s and SOC λ we observe three regions that exhibit BFSs [phases II, III, and IV in Fig. 2(a)]. Phase II has two BFSs, while phases III and IV have four BFSs. We note that phases III and IV are separated by a thin line, where the two pairs of BFSs touch each other, cf. dashed line in Fig. 2(a) and Figs. S1(b) and S1(e) in the SM. In the next section we discuss the conductance and noise spectra of phases I, II, and III, and the transitions between them. The transition between phases III and IV is investigated in the SM [31].

Differential conductance and noise spectroscopy. Let us now turn to the transport properties, which we use to characterize the different phases I, II, and III of Fig. 2. In order to compute the differential conductance G and the current-noise ratio F , we impose periodic boundary conditions in the y direction of the considered device (Fig. 1). For concreteness we fix the value of the Rashba SOC to $\lambda/\mu_n = 1.0$ and study the transport properties as a function of Zeeman field V_s , with $V_n = 5V_s$. By examining Fig. 2, we find that for $\lambda/\mu_n = 1.0$, phase I (no BFSs) occurs for $V_s/\Delta < 0.1405$, phase II (two BFSs) is realized for $0.1405 < V_s/\Delta < 0.2388$, while phase III (four BFSs) exists for $V_s/\Delta > 0.2388$. These three phases are marked in Figs. 3 and 4 by white, green, and gray color, respectively.

We first examine the differential conductance at zero temperature, which we compute within the Blonder–Tinkham–Klapwijk formalism (BTK) [33]. For this purpose we apply a bias voltage eV to our device, which injects electrons from $i_x = -\infty$. The differential conductance is evaluated by

$$G(eV) = \frac{e^2}{h} \sum_{k_y} \text{Tr}[\hat{\sigma}_0 - \hat{R}_e + \hat{R}_h]_{E=eV}, \quad (6a)$$

$$\hat{R}_\alpha = \hat{r}_\alpha \hat{r}_\alpha^\dagger, \quad (\hat{r}_\alpha)_{\sigma,\sigma'} = r_{\sigma\sigma'}^{\alpha e}(k_y, E), \quad (6b)$$

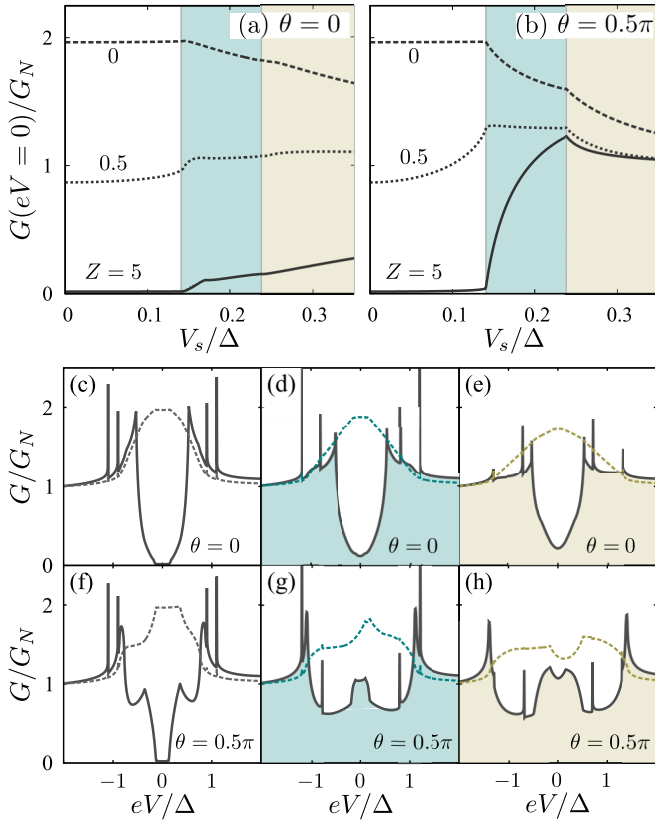


FIG. 3. Panels (a) and (b) show the normalized zero-bias differential conductance G/G_N as a function of Zeeman field V_s for the field angles $\theta = 0$ and $\theta = \pi/2$, respectively, with three different junction transparencies Z . Panels (c)–(h) display the differential conductance as a function of bias voltage for field angle $\theta = 0$ (top row), $\theta = \pi/2$ (bottom row), junction transparency $Z = 0$ (dashed lines), and $Z = 5$ solid lines. The first column [(c) and (f)] corresponds to phase I (no BFSs) with $V_s = 0.1\Delta$, the second column [(d) and (g)] represents phase II (two BFSs) with $V_s = 0.2\Delta$, while the third column [(e) and (h)] is in phase III (four BFSs) with $V_s = 0.3\Delta$.

where $r_{\sigma\sigma'}^{\alpha e}(k_y, E)$ for $\alpha = e$ ($\alpha = h$) represents the normal (Andreev) reflection coefficient from an electron with spin σ' to an electron (hole) with spin σ at momentum k_y and energy E . $\sum'_{k_y}$ denotes the summation over k_y of the propagating channels. The reflection coefficients are calculated using lattice Green's function techniques [29,30]. In the figures we show the normalized differential conductance G/G_N , which is obtained by dividing by the normal conductance G_N calculated by setting $\Delta = eV = V_s = V_n = 0$.

In Figs. 3(c)–3(h) we show G/G_N as a function of bias voltage for the three phases of Fig. 2(a) and the field angles $\theta = 0$ [top row, panels (c)–(e)] and $\theta = \pi/2$ [bottom row, panels (f)–(h)]. Figures 3(a) and 3(b) display G/G_N at zero bias as a function of Zeeman field V_s for $\theta = 0$ and $\theta = \pi/2$, respectively. Let us first focus on the tunneling limit with low junction transparency $Z = 5$ (solid lines). For the tunneling limit we observe in all panels (c)–(h) two peaks symmetrically located around $\pm\Delta$. These double peaks are the coherence peaks of the parent superconductor, which are split by the applied Zeeman field (cf. Fig. S2 of the SM [31]). In addition there are smaller peaks (mostly at smaller bias),

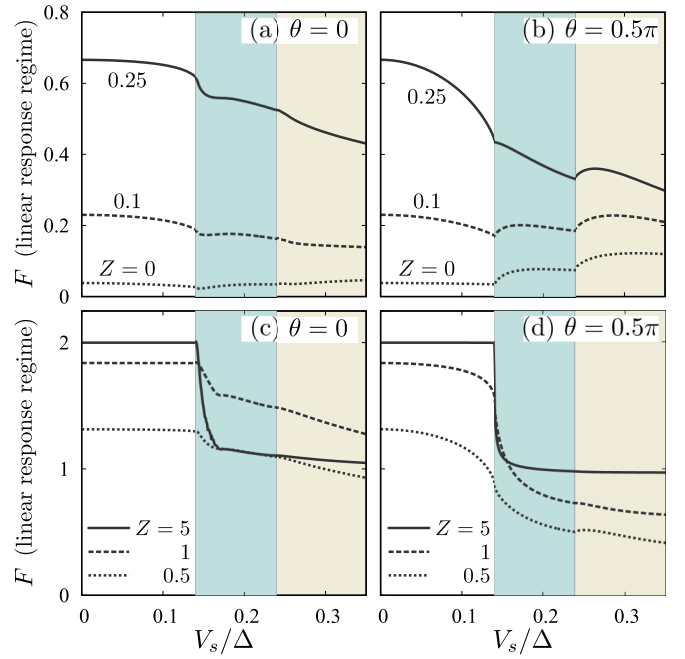


FIG. 4. Fano factor F within the linear response regime, Eq. (8), as a function of Zeeman field V_s for six different junction transparencies Z . Panels (a) and (c) correspond to field angle $\theta = 0$, while in panels (b) and (d) the field angle is set to $\theta = \pi/2$. The white, green, and gray regions correspond to phase I (no BFSs), phase II (two BFSs), and phase III (four BFSs), respectively, see Fig. 2(a).

which originate from the proximity induced superconductivity of the semimetal. Around zero-bias voltage we observe clear signatures of the BFSs, i.e., in the tunneling limit there is residual conductance at zero bias originating from the finite zero-energy density of states of the BFSs. The conductance around zero bias increases with increasing V_s and shows kink singularities at the phase transition points, where the number of BFSs jumps from zero to two and from two to four (see solid lines in Figs. 3(a) and 3(b)). Interestingly, the zero-bias conductance is larger for $\theta = \pi/2$ than for $\theta = 0$. This is due to the fact that for $\theta = 0$ most of the incident electrons (i.e., those with $|k_y/\pi| \lesssim 0.05$) cannot couple to the BFSs, while for $\theta = \pi/2$ almost all incident electrons can tunnel into the BFSs [see Figs. 1(b)–1(d)].

Also in the ballistic limit with a completely transparent junction ($Z = 0$), distinct signatures of the BFSs can be observed in G/G_N around zero bias (dashed lines in Fig. 3). That is, in the absence of BFSs all incident electrons show perfect Andreev reflection yielding $G/G_N = 2$, while in the presence of BFSs only parts of the electrons Andreev reflect leading to a suppression of G/G_N , with $1 < G/G_N < 2$. As in the tunneling limit, we observe distinct kink singularities at the transition points that separate phases with different numbers of BFSs [see dashed lines in Figs. 3(a) and 3(b)]. In passing, we note that the conductance spectra (particularly in the ballistic limit $Z = 0$) are not symmetric around zero bias. The reason for this is the broken time-reversal and broken inversion symmetry, as discussed in more detail in the SM [31].

Let us now turn to the noise spectra of the BFSs. For this purpose we compute the current-noise ratio (Fano factor)

within the linear response regime, i.e., for $eV \ll \Delta$, where the BFSs show a dominant contribution. The current-noise ratio is given by the ratio $F = S/e\bar{I}$ between the zero-frequency noise power

$$S = 2 \int_{-\infty}^{\infty} \overline{\delta I(0)\delta I(t)} dt \quad (7)$$

and the time averaged current at zero temperature $\bar{I}$. Here, $\delta I(t) = I(t) - \bar{I}$ denotes the deviation of the current at time t from the averaged current $\bar{I}$. Within the linear response regime, the Fano factor at zero bias simplifies to

$$F = \frac{P(eV = 0)}{eG(eV = 0)}, \quad (8)$$

where the noise power $P(eV = 0)$ [22,23] is evaluated within the BTK formalism [33] as

$$P(0) = \frac{e^3}{h} \sum_{k_y} \text{Tr} \left[\sum_{\alpha=e,h} \hat{R}_{\alpha} (\hat{\sigma}_0 - \hat{R}_{\alpha}) + 2\hat{R}_e \hat{R}_h \right]_{E=0}. \quad (9)$$

As before, we compute the reflection coefficients entering in the above formula by use of lattice Greens function techniques [29,30].

In Fig. 4 we present the zero-bias Fano factor, Eq. (8), as a function of Zeeman field V_s for six different junction transparencies Z . The left column [(a) and (c)] corresponds to the field angle $\theta = 0$, while the right column [(b) and (d)] shows the results for the field angle $\theta = \pi/2$. The different phases of Fig. 2(a), with different numbers of BFSs, are indicated by the white, green, and gray colors, respectively. In the tunneling limit with low junction transparency $Z = 5$ the Fano factor gives us information about the fundamental charge of the quasiparticles that tunnel across the junction. In phase I without BFSs (white regions), the Fano factor takes the value of two, since the tunneling current is carried exclusively by Cooper pairs composed of two electrons. In phases II and III, however, which exhibit a large number of BQPs forming BFSs (green and gray regions), the Fano factor decreases to a value substantially lower than two. This indicates that the current is carried by both Cooper pairs and BQPs, which have an effective charge substantially smaller than that of Cooper pairs. Remarkably, if we set the field angle to $\theta = \pi/2$, such that almost all incident electrons can couple to the BFSs, the Fano factor drops below one [solid line in Fig. 4(d)], implying that almost all of the current is carried by the BQPs.

In the ballistic limit with transparent junction $Z = 0$, we can obtain information about the types of scattering processes that occur at the junction barrier. In phase I without BFSs we observe that the Fano factor is almost zero [dotted lines in Figs. 4(a) and 4(b)], as almost all incident electrons undergo Andreev reflection such that the noise power P is almost zero. In phases II and III with BFSs, however, the Fano factor increases, indicating that some electrons are transformed into

BQPs of the BFSs, while others Andreev reflect. By setting the field angle to $\theta = \pi/2$, we can maximize the coupling between the lead electrons and the BFS, such that more electrons transmit into the BFSs, thereby increasing the Fano factor [cf. dotted lines in Figs. 4(a) and 4(b)]. As in Fig. 3 we observe distinct kinks in the Fano factor at the transition points that separate the different phases.

Conclusions. In summary, we have investigated the differential conductance and noise spectroscopy of a NS junction with BFSs. We find that the current-noise ratio (Fano factor) provides key fingerprints of the presence of BFSs (see Fig. 4): In the tunneling limit the Fano factor drops from two to about one, whenever there are BFSs present. In the ballistic limit, on the other hand, the Fano factor jumps from zero to a finite value due to the presence of BFSs. Moreover, the Fano factor as a function of Zeeman field exhibits kink discontinuities at the transition lines that separate phases with different numbers of BFSs. We hope that these predictions will help experimentalists to identify the presence of BFSs in their superconducting hybrid devices.

We note that the creation of BFSs in semiconductor/superconductor hybrid devices requires Zeeman fields larger than the proximity-induced pair potential in the semiconductor, but smaller than the critical field of the parent superconductor. This condition is realized for hybrid devices, such as Al/InAs heterostructures [16,24], where the g factor of the semiconductor is large, while the proximity-induced pair potential is small compared to the parent superconductor. While we have focused in this Letter on superconducting hybrids with BFSs, the obtained results are also of relevance to other systems with BFSs, for example, multiband superconductors with broken time-reversal symmetry [6–9].

In closing, we mention several promising directions for future research. First, it would be interesting to investigate other types of junctions, e.g., NSN and SNS junctions, involving superconductors with BFSs, and to determine the signatures of BFSs in the Josephson current. Experimentally, it should be straightforward to fabricate NSN and SNS junctions using existing thin film technology. Second, it would be of interest to study thermal transport properties of BFSs. While conventional superconductors are perfect thermal insulators at low temperature, superconductors with BFSs are expected to show substantial thermal transport carried by the large density of BQPs near zero energy. Third, the interplay between BFSs and Andreev surface bound states (or interface bound states) and its effects on the transport properties would be a worthwhile topic for future research.

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