Surface Bound States and Spin Currents in Noncentrosymmetric Superconductors

A. B. Vorontsov,[1,*] I. Vekhter,[1] and M. Eschrig[2]

[1]*Department of Physics and Astronomy, Louisiana State University, Baton Rouge, Louisiana, 70803, USA*
[2]*Institut für Theoretische Festkörperphysik and DFG-Center for Functional Nanostructures, Universität Karlsruhe, D-76128 Karlsruhe, Germany*

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We investigate the ground state properties of a noncentrosymmetric superconductor near a surface. We determine the spectrum of Andreev bound states due to surface-induced mixing of bands with opposite spin helicities for a Rashba-type spin-orbit coupling. We find that the order parameter suppression qualitatively changes the bound state spectrum. The spin structure of Andreev states leads to a spin supercurrent along the interface, which is strongly enhanced compared to the normal state spin current. Particle and hole coherence amplitudes show Faraday-like rotations of the spin along quasiparticle trajectories.

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The role of chirality and spin-orbit coupling in materials and nanostructures is a very active and promising subject in the fields of spintronics, superconductivity, and magnetism [1–3]. The unusual properties of noncentrosymmetric (NCS) materials originate from the crystal structure that lacks a center of inversion, allowing for pronounced spin-orbit (SO) coupling that is odd in the electron momentum, and leading to a chiral ground state. The resulting two-band nature of NCS metals leads to effects reminiscent of semiconductor physics, such as birefringence and spin polarization of the electron wave packet [4]. Especially promising is the presence of charge-neutral spin currents in the ground state [5,6].

Since understanding of interface physics is one of the foundations for all potential applications, it is of pivotal interest to investigate how the physical properties of NCS materials are modified near surfaces. The key point here is that interface scattering in materials with strong spin-orbit effects is spin active. Spins dominate the surface physics, and any successful theoretical treatment must take this into account.

The recently discovered class of NCS superconductors [7–10] combines the strong SO coupling that governs the metallic bands with a nontrivial spin structure of the superconducting (SC) order parameter [11–13] due to lack of parity. As a result, one may expect that spin transport in the SC phase exhibits novel features compared to usual superconductors. These features are especially prominent near surfaces and interfaces, where the physics is controlled by the Andreev bound states (ABS), built as a result of particle-hole coherent scattering. ABS are crucially important in unconventional superconductors [14], where the phase variation of the order parameter (OP) on the Fermi surface [15] and the pairbreaking near interfaces lead to a midgap peak in the density of states (DOS) at the surface. The ABS states control thermodynamic properties and stability of the surface phases [16,17], and govern transport across interfaces [18–20].

In this Letter we study the Andreev states and spin currents at the surface of a NCS superconductor. We show that due to the nontrivial spin structure of the superconducting order, (i) a finite-bias peak appears in the surface DOS, and (ii) Andreev states acquire spin polarization and carry a spin supercurrent. Spin currents are enhanced near the surface already in the normal state, but the spin current in the SC phase has a significantly greater magnitude. We develop a detailed theory of these effects, and compare them with the previously studied spin currents in metallic and semiconducting systems.

For a noncentrosymmetric material it is convenient to perform a canonical transformation from a spin basis (with fermion annihilation operators $c_{\mathbf{k}\mu}$ for spin $\mu = \uparrow, \downarrow$) to the so-called helicity basis ($b_{\mathbf{k}s}$ with helicity $s = \pm$), that diagonalizes the kinetic part of the Hamiltonian,

$$\mathcal{H}_{\rm kin} = \sum_{\mathbf{k}\mu\nu} c^\dagger_{\mathbf{k}\mu}(\xi_{\mathbf{k}} + \alpha \mathbf{g}_{\mathbf{k}}\boldsymbol{\sigma})_{\mu\nu} c_{\mathbf{k}\nu} = \sum_{\mathbf{k}s} \varepsilon_{\mathbf{k}s} b^\dagger_{\mathbf{k}s} b_{\mathbf{k}s}. \quad (1)$$

Here, $\xi_{\mathbf{k}}$ is the band dispersion relative to the chemical potential in the absence of SO interaction, α is the SO coupling strength, $\boldsymbol{\sigma}$ is the vector of Pauli matrices, and $\mathbf{g}_{\mathbf{k}} = -\mathbf{g}_{-\mathbf{k}}$ is a normalized (see below) SO vector [11,12], Fig. 1. The helicity band dispersion is $\varepsilon_{\mathbf{k}\pm} = \xi_{\mathbf{k}} \pm \alpha |\mathbf{g}_{\mathbf{k}}|$. In each helicity band the orientation of the quasiparticle spin is fixed with respect to its momentum.

The Hamiltonian (1) is time reversal invariant but lifts the spin degeneracy. The transformation from spin to helicity basis, $U_{\mathbf{k}}$, is defined by $U_{\mathbf{k}}(\mathbf{g}_{\mathbf{k}}\boldsymbol{\sigma})U^\dagger_{\mathbf{k}} = |\mathbf{g}_{\mathbf{k}}|\sigma_3$, and determined by the direction of the $\mathbf{g}$ vector in $\mathbf{k}$ space,

$$U_{\mathbf{k}} = e^{-i(\theta_{\mathbf{g}}/2)\mathbf{n}_{\mathbf{g}}\boldsymbol{\sigma}}, \qquad \mathbf{n}_{\mathbf{g}} = \mathbf{g}_{\mathbf{k}} \times \hat{\mathbf{z}}/|\mathbf{g}_{\mathbf{k}} \times \hat{\mathbf{z}}|, \quad (2)$$

where $\hat{\mathbf{z}}$ is the unit vector in the z direction, and $\theta_{\mathbf{g}}$ is the polar angle between $\mathbf{g}_{\mathbf{k}}$ and $\hat{\mathbf{z}}$ [12].

To describe superconductivity we use the Nambu-Gor'kov formalism modified for a helical basis. We define the helical counterpart $\hat{B}^\dagger_{\mathbf{k}} = (b^\dagger_{\mathbf{k}+}, b^\dagger_{\mathbf{k}-}, b_{\mathbf{k}+}, b_{\mathbf{k}-})$ to the

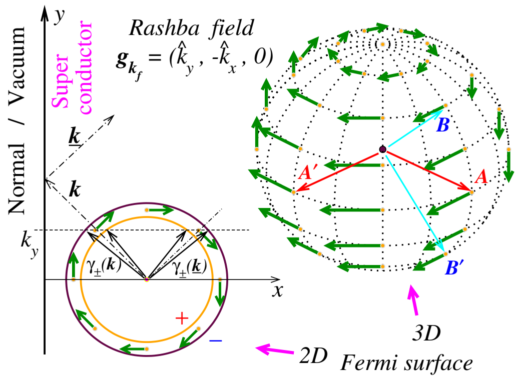

FIG. 1 (color online). A map of the spin-orbit vector in momentum space for the Rashba form $\mathbf{g_k} = \hat{\mathbf{k}} \times \hat{\mathbf{z}}$. On reflection the spin-orbit vector $\mathbf{g_k}$ may change, e.g., from $A \rightarrow A'$, or not, $B \rightarrow B'$. The scattering geometry is shown on the left.

Nambu spinor $\hat{C}^\dagger_{\mathbf{k}} = (c^\dagger_{\mathbf{k}\uparrow}, c^\dagger_{\mathbf{k}\downarrow}, c_{\mathbf{k}\uparrow}, c_{\mathbf{k}\downarrow})$ by $\hat{B}_{\mathbf{k}} = \hat{U}_{\mathbf{k}}\hat{C}_{\mathbf{k}}$, $\hat{U}_{\mathbf{k}} = \mathrm{diag}(U_{\mathbf{k}}, U^*_{\mathbf{k}})$ and construct 4×4 retarded Green's functions in the helicity basis, $\hat{G}_{\mathbf{k}_1\mathbf{k}_2}(t_1, t_2) = -i\theta(t_1 - t_2)\langle\{\hat{B}_{\mathbf{k}_1}(t_1), \hat{B}^\dagger_{\mathbf{k}_2}(t_2)\}\rangle_{\mathcal{H}}$, where $\hat{B}(t)$ are Heisenberg operators, the braces denote an anticommutator, $\langle\ldots\rangle_{\mathcal{H}}$ is a grand canonical average, and θ is the usual step function.

We employ the quasiclassical method [21] for treating the inhomogeneous surface problem. In the materials of interest $\alpha|\mathbf{g}_{\mathbf{k}_f}| \ll E_f$ for any Fermi momentum $\mathbf{k}_f$, where E_f is the Fermi energy. In addition, the superconducting energy scales (transition temperature T_c and the gap Δ) are much smaller than E_f. Under these conditions quasiparticles with different helicity can be assigned to a common Fermi surface and propagate coherently along a common classical trajectory over distances much longer than the Fermi wavelength. We normalize $\mathbf{g_k}$, $\langle \mathbf{g}^2_{\mathbf{k}_f}\rangle = 1$, where $\langle\ldots\rangle$ denotes a Fermi surface average. The quasiclassical propagator is then obtained as $\hat{g}(\mathbf{k}_f, \mathbf{R}, \epsilon, t) = \hat{\tau}_3 \int d\xi_{\mathbf{k}} \times \int (d\mathbf{q})(d\tau) e^{i(\mathbf{qR}+\epsilon\tau)}\hat{G}_{\mathbf{k}+\mathbf{q}/2,\mathbf{k}-\mathbf{q}/2}(t + \frac{\tau}{2}, t - \frac{\tau}{2})$ where $\hat{\tau}_3$ is the Pauli matrix in the particle-hole space. Using $U_{-\mathbf{k}}U^\dagger_{\mathbf{k}} = i\mathbf{n_g}\boldsymbol{\sigma}$ and the fermionic anticommutation relations for the b and $b^\dagger$, we derive the fundamental symmetry relations for the 2×2 Nambu matrix components, $g(\epsilon, \mathbf{k}_f)_{22} = [(\mathbf{n_g}\boldsymbol{\sigma})g(-\epsilon, -\mathbf{k}_f)_{11}(\mathbf{n_g}\boldsymbol{\sigma})]^*$ and $g(\epsilon, \mathbf{k}_f)_{21} = [(\mathbf{n_g}\boldsymbol{\sigma})g(-\epsilon, -\mathbf{k}_f)_{12}(-i\sigma_2)(\mathbf{n_g}\boldsymbol{\sigma})]^*(i\sigma_2)$.

Standard procedure [21] yields the Eilenberger equation in the helicity basis,

$$[\varepsilon\hat{\tau}_3 - \alpha\hat{v}_{\mathrm{SO}} - \hat{\Delta}, \hat{g}] + i\mathbf{v}_f\boldsymbol{\nabla}\hat{g} = \hat{0} \tag{3}$$

with normalization $\hat{g}^2 = -\pi^2\hat{1}$. Here, ε is the energy, $\hat{v}_{\mathrm{SO}} = |\mathbf{g}_{\mathbf{k}f}|\sigma_3\hat{\tau}_3$, and $\hat{\Delta}$ is the superconducting OP. The velocity renormalization of order $\alpha/E_f \ll 1$ is neglected. We choose a separable pairing interaction consistent with the form of the gap, and determine $\hat{\Delta}$ self consistently with $\hat{g}$. In NCS superconductors the OP is a mixture of spin singlet (Δ_s) and triplet (Δ_t) components [13,22]. In the real gauge it is given by

$$\hat{\Delta} = \mathcal{Y}(\mathbf{k}_f)[\Delta_s(\mathbf{R})\hat{1} + \Delta_t(\mathbf{R})\hat{v}_{\mathrm{SO}}(\mathbf{k}_f)](i\sigma_2)\hat{\tau}_1, \tag{4}$$

where the basis function $\mathcal{Y}(\mathbf{k}_f)$ transforms according to one of the irreducible representations of the crystal point group, and $\langle\mathcal{Y}^2(\mathbf{k}_f)\rangle = 1$. With the gap functions in the helicity bands, $\Delta_\pm = \Delta_s \pm \Delta_t|\mathbf{g}_{\mathbf{k}_f}|$, the order parameter is $\Delta = \{\Delta_+\sigma^+ - \Delta_-\sigma^-\}\mathcal{Y}$, where $\sigma^\pm = (\sigma_1 \pm i\sigma_2)/2$.

We parametrize the Green's function by the coherence functions for particles and holes, γ and $\tilde{\gamma}$ (2×2 spin matrices), which allow a very intuitive physical interpretation of the Andreev scattering processes [23],

$$\hat{g} = -i\pi\begin{pmatrix} 1-\gamma\tilde{\gamma} & 0 \\ 0 & 1-\tilde{\gamma}\gamma \end{pmatrix}^{-1}\begin{pmatrix} 1+\gamma\tilde{\gamma} & 2\gamma \\ -2\tilde{\gamma} & -1+\tilde{\gamma}\gamma \end{pmatrix}. \tag{5}$$

Fundamental symmetry relates γ and $\tilde{\gamma}$ in the helicity basis by $\tilde{\gamma}(\epsilon, \mathbf{k}_f) = [(\mathbf{n_g}\boldsymbol{\sigma})\gamma(-\epsilon, -\mathbf{k}_f)(-i\sigma_2)(\mathbf{n_g}\boldsymbol{\sigma})]^*(i\sigma_2)$. In the bulk, $\gamma = \gamma^0_+\sigma^+ - \gamma^0_-\sigma^-$, and $\tilde{\gamma} = \tilde{\gamma}^0_-\sigma^+ - \tilde{\gamma}^0_+\sigma^-$, with $\gamma^0_\pm(\epsilon, \mathbf{k}_f) = -\Delta_\pm(\mathbf{k}_f)/(\epsilon + i\sqrt{|\Delta_\pm(\mathbf{k}_f)|^2 - \epsilon^2})$ and $\tilde{\gamma}^0_\pm(\mathbf{k}_f, \epsilon) = \gamma^0_\pm(-\mathbf{k}_f, -\epsilon)^*$.

The surface bound states are determined by the poles of the Green's function, Eq. (5). We consider specular reflection, whereby the component of $\mathbf{k}$ normal to the surface changes sign, $\mathbf{k} \rightarrow \underline{\mathbf{k}}$; see Fig. 1. We find $\gamma_{\mathbf{k}}$ ($\tilde{\gamma}_{\underline{\mathbf{k}}}$), by integrating forward (backward) along the incoming, $\mathbf{k}$, (outgoing, $\underline{\mathbf{k}}$) trajectory starting from the values in the bulk [23]. In contrast, $\gamma_{\underline{\mathbf{k}}}$ and $\tilde{\gamma}_{\mathbf{k}}$, are determined from the boundary conditions at the surface. Since the interface is nonmagnetic, $\hat{g}$ in the spin basis, $\hat{g}^s_{\mathbf{k}} = \hat{U}^\dagger_{\mathbf{k}}\hat{g}_{\mathbf{k}}\hat{U}_{\mathbf{k}}$, is continuous at the surface. This leads to a surface-induced mixing of the helicity bands according to $U^\dagger_{\underline{\mathbf{k}}}\gamma_{\underline{\mathbf{k}}}U^*_{\underline{\mathbf{k}}} = \gamma^s_{\underline{\mathbf{k}}} = \gamma^s_{\mathbf{k}} = U^\dagger_{\mathbf{k}}\gamma_{\mathbf{k}}U^*_{\mathbf{k}}$ and $U^T_{\mathbf{k}}\tilde{\gamma}_{\mathbf{k}}U_{\mathbf{k}} = \tilde{\gamma}^s_{\mathbf{k}} = \tilde{\gamma}^s_{\underline{\mathbf{k}}} = U^T_{\underline{\mathbf{k}}}\tilde{\gamma}_{\underline{\mathbf{k}}}U_{\underline{\mathbf{k}}}$. From Eq. (5), the bound states correspond to the zero eigenvalues of the matrix $1 - \gamma_{\mathbf{k}}\tilde{\gamma}_{\mathbf{k}} = 1 - \gamma_{\mathbf{k}}U^*_{\mathbf{k}}U^T_{\underline{\mathbf{k}}}\tilde{\gamma}_{\underline{\mathbf{k}}}U_{\underline{\mathbf{k}}}U^\dagger_{\mathbf{k}}$ at the surface, and we derive our final equation for the ABS energies via the surface amplitudes in the helicity basis

$$(1 + \gamma_+\tilde{\gamma}_+)(1 + \gamma_-\tilde{\gamma}_-) = -(1 + \gamma_+\tilde{\gamma}_-)(1 + \gamma_-\tilde{\gamma}_+)\mathcal{M}. \tag{6}$$

The "mixing" factor $\mathcal{M}$ is determined by the change of $\mathbf{g_k}$ under reflection $\mathbf{k} \rightarrow \underline{\mathbf{k}}$ at the surface,

$$\mathcal{M} = \frac{\sin^2\frac{\theta_{\mathbf{g}} - \theta_{\underline{\mathbf{g}}}}{2} + \sin^2\frac{\theta_{\mathbf{g}} + \theta_{\underline{\mathbf{g}}}}{2}\tan^2\frac{\phi_{\mathbf{g}} - \phi_{\underline{\mathbf{g}}}}{2}}{\cos^2\frac{\theta_{\mathbf{g}} - \theta_{\underline{\mathbf{g}}}}{2} + \cos^2\frac{\theta_{\mathbf{g}} + \theta_{\underline{\mathbf{g}}}}{2}\tan^2\frac{\phi_{\mathbf{g}} - \phi_{\underline{\mathbf{g}}}}{2}}, \tag{7}$$

where $\theta_{\mathbf{g}}$, $\phi_{\mathbf{g}}$ and $\theta_{\underline{\mathbf{g}}}$, $\phi_{\underline{\mathbf{g}}}$ are the polar and azimuthal angles of $\mathbf{g_k}$ and $\mathbf{g}_{\underline{\mathbf{k}}}$, respectively. If $\mathbf{g_k} = \mathbf{g}_{\underline{\mathbf{k}}}$ ($B \rightarrow B'$ in Fig. 1) there is no helicity band mixing, $\mathcal{M} = 0$, and we recover the conditions for ABS in superconductors with no SO coupling. The limit $\mathcal{M} \rightarrow \infty$ describes pure interband scattering. For a general Fermi surface and an arbitrary vector $\mathbf{g_k}$ ($A \rightarrow A'$ in Fig. 1) a finite $\mathcal{M}$ determines the relative weights of intraband and interband scattering, and Eq. (6) gives the bound state energy.

We show now that the suppression of the anisotropic (triplet) component of the OP in Eq. (4) near the surface drastically modifies the ABS spectrum, and the surface DOS, $N(\varepsilon, \mathbf{k}_f) = -\frac{N_f}{2\pi} \mathrm{Im}\,\mathrm{Tr}\{g(\varepsilon, \mathbf{k}_f)\}$, where Tr is a 2×2 spin trace, and N_f is the normal state DOS. The salient features are clear from considering a 2D material with the Rashba-type SO coupling $\alpha = \alpha_R k_f$, $\mathbf{g}_\mathbf{k} = (\mathbf{k} \times \hat{\mathbf{z}})/k_f = (k_y, -k_x, 0)/k_f$, and a triplet order parameter, $\Delta_+ = -\Delta_- = \Delta$; we find only quantitative differences for 3D Fermi surface, $\Delta_+ \neq -\Delta_-$ and other $\mathbf{g}_\mathbf{k}$'s.

To obtain insight in the role of the OP suppression, we consider first a simple model where $\Delta = 0$ in a layer of width W next to the surface; see Fig. 2. Trajectories incident at an angle ϕ travel through a non-SC region of an effective width $2D = 2W/\cos\phi$. In this case Eq. (7) reads $\mathcal{M} = \cot^2\phi$, the surface coherence amplitudes gain a phase factor, $\gamma_\pm = \gamma^0_\pm e^{i2\varepsilon D/v_f}$, $\tilde{\gamma}_\pm = \tilde{\gamma}^0_\pm e^{i2\varepsilon D/v_f}$, Eq. (6) yields $\mathrm{Im}^2(\tilde{\gamma}^0_+ e^{i2\varepsilon D/v_f}) = \mathrm{Re}^2(\tilde{\gamma}^0_+ e^{i2\varepsilon D/v_f})\mathcal{M}$, and the bound states are given by

$$\varepsilon/\sqrt{\Delta_0^2 - \varepsilon^2} = -\tan(2W\varepsilon/v_f \cos\phi \pm \phi). \qquad (8)$$

Solutions of this equation are shown in Fig. 2. The "principal" modes with energies away from the continuum edge contribute the most to the subgap DOS. $W = 0$ (no gap suppression) reproduces the result of Ref. [24]: each incoming trajectory yields a bound state at a different energy. For $W \neq 0$ the main mode $\varepsilon_{\mathrm{bs}}(\phi)$ develops a maximum at $\varepsilon^\star < \Delta_0$, and we expect a peak in the surface DOS near $\varepsilon^\star$ due to abundance of trajectories contributing to $N(\varepsilon^\star)$.

The fully self-consistent solution, shown in Fig. 3(b), confirms this. Note that $\Delta \neq 0$ at the surface, Fig. 3(a), as in other unconventional superconductors misaligned with respect to the interface [16]. Crucially, self-consistency does yield a strong peak in the surface DOS below the gap at a finite energy, in qualitative contrast to earlier results [24,25]. Experimentally accessing this peak by point contact tunneling requires a sufficiently wide tunneling cone as the feature arises from the trajectories at intermediate incident angles; see Fig. 2.

These ABS have unusual spin structure. Figures 3(c) and 3(d) show the spin-resolved density of states, $N_{\uparrow\downarrow} = N \pm N^Z$, where N is the net DOS and $N^\alpha(\epsilon, \mathbf{k}_f, \mathbf{x}) = -\frac{N_f}{2\pi} \mathrm{Im}\,\mathrm{Tr}\{\sigma^\alpha g(\epsilon, \mathbf{k}_f, \mathbf{x})\}$. At the interface $N^X = N^Y = 0$. The states corresponding to different branches of Eq. (8) have opposite spin polarization. Since the spin polarization changes sign for reversed trajectories, the Andreev states carry spin current along the interface.

Spin currents exist in NCS materials because the spin is not conserved, and consequently precession terms enter the continuity equation, $\partial_t S^\alpha(\mathbf{x}) + \boldsymbol{\nabla} \cdot \boldsymbol{\Pi}^\alpha(\mathbf{x}) = P^\alpha(\mathbf{x})$ [5]. Here, the spin density, $S^\alpha(\mathbf{x}) = \frac{1}{2}\,\mathrm{Tr} \int d\mathbf{k}\, \sigma^\alpha G(\mathbf{k}, \mathbf{x})$, the spin current, $\boldsymbol{\Pi}^\alpha(\mathbf{x}) = \frac{1}{4}\,\mathrm{Tr} \int d\mathbf{k} \{\sigma^\alpha, \mathbf{v}_\mathbf{k}\} G(\mathbf{k}, \mathbf{x})$, and the precession $P^\alpha(\mathbf{x}) = \frac{1}{2i}\,\mathrm{Tr} \int d\mathbf{k} [\sigma^\alpha, \mathbf{v}_\mathbf{k} \cdot \mathbf{k}] G(\mathbf{k}, \mathbf{x})$, (where $[\bullet, \bullet]$ is a commutator, and $\mathbf{v}_\mathbf{k} = \mathbf{k}_f/m + \alpha_R[\hat{\mathbf{z}} \times \boldsymbol{\sigma}]$ is the band velocity), are all given in terms of Green's functions at imaginary relative time $\tau = -i0$. For the Rashba case, the precession terms are related to spin currents via the relations $P^X = -2m\alpha_R \Pi^Z_x$, $P^Y = -2m\alpha_R \Pi^Z_y$, $P^Z = 2m\alpha_R(\Pi^X_x + \Pi^Y_y)$ [26].

We first consider the spin currents in the normal state. The bulk value, $\Pi^Y_x = -\Pi^X_y = \Pi^{\mathrm{bulk}}_{\mathrm{spin}} = m^2\alpha_R^3/3\pi$ agrees with Ref. [5]. To determine the surface spin currents we find the Green's function for a surface modeled as a δ-function barrier at $x = 0$ of strength U. The Dyson equation in 2×2 spin space reads $G^{-1} = [G^{(0)}]^{-1} - U\delta(x)$, where $[G^{(0)}_\mathbf{k}]^{-1} = \varepsilon - \xi_\mathbf{k} - \alpha_R(\mathbf{k} \times \hat{\mathbf{z}})\boldsymbol{\sigma}$. For an impenetrable surface ($U \to \infty$) the solution is (for fixed k_y) [27]

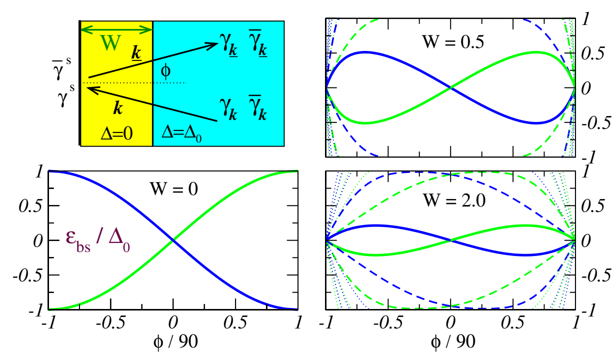

FIG. 2 (color online). Bound state energy as a function of the impact angle for different widths W (in units of $v_f/2\Delta_0$) of the order parameter suppression region. Dark (light) lines correspond to the plus (minus) sign in Eq. (8). Solid lines: principal mode; broken lines: higher multiple reflection modes.

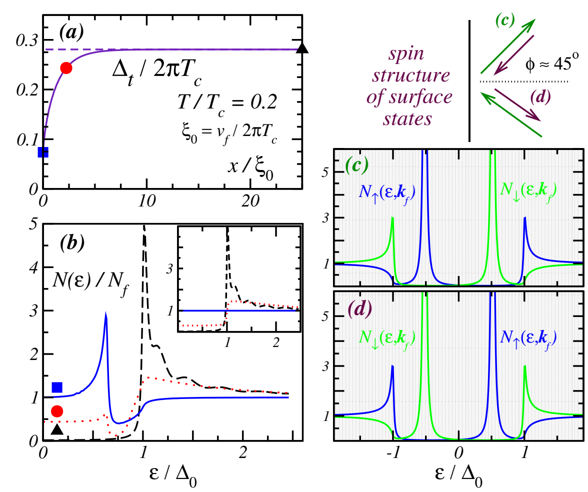

FIG. 3 (color online). Structure of the surface states for a Rashba triplet superconductor ($\Delta_+ = -\Delta_-$). (a) Order parameter suppression; (b) DOS at locations indicated in (a); the large subgap peaks are due to the suppression of Δ. Inset: DOS for a homogeneous order parameter; note the absence of any subgap features. (c, d) Spin-resolved surface DOS for two trajectories; $N_\uparrow(\epsilon, \mathbf{k}_f)$ (dark) and $N_\downarrow(\epsilon, \mathbf{k}_f)$ (light) correspond to dark and light branches in Fig. 2.

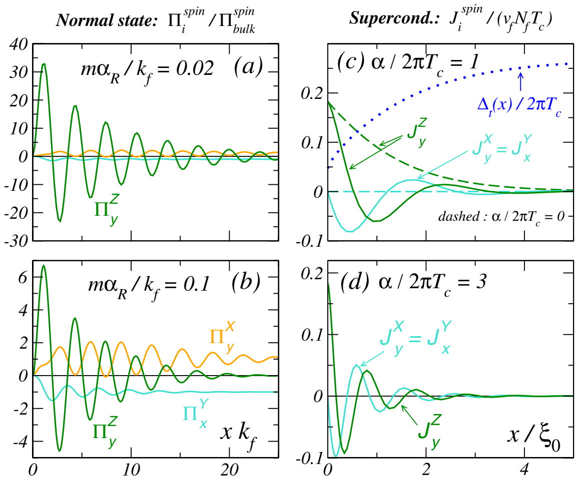

FIG. 4 (color online). Spin currents near the surface. Left: Rashba metal. The nonvanishing components in the bulk are $\Pi_x^Y = -\Pi_y^X$. Near the surface Π_y^Z is large; see text. Right: Spin currents and the order parameter in the SC state.

$$G_{k_x k_x'} = G^{(0)}_{k_x} 2\pi\delta(k_x - k_x') - G^{(0)}_{k_x} \frac{1}{\int \frac{dp_x}{2\pi} G^{(0)}_{p_x}} G^{(0)}_{k_x'}. \quad (9)$$

We solve Eq. (9) numerically, and show the normal state surface spin currents in Figs. 4(a) and 4(b). The most prominent new feature is a large surface current Π_y^Z with out of plane spin polarization [28] that flows along the surface, and decays rapidly into the bulk on the scale similar to that of Friedel oscillations. This component is related to Π_x^Y via the continuity equation, $\Pi_y^Z(x) = -1/(2m\alpha_R) d\Pi_x^Y(x)/dx$. As a result, this component is much greater, by a factor of order $k_f/m\alpha_R$, than the bulk spin currents in the normal state.

The SC spin current, shown in Figs. 4(c) and 4(d), is defined in the quasiclassical method relative to the normal state,

$$\mathbf{J}^{\alpha} \equiv \mathbf{\Pi}^{\alpha} - \mathbf{\Pi}^{\alpha}_N = \int_{-\infty}^{\infty} d\epsilon n_f(\epsilon) \langle \mathbf{v}_f N^{\alpha}(\epsilon, \mathbf{k}_f, \mathbf{x}) \rangle, \quad (10)$$

where $n_f(\epsilon)$ is the Fermi function. The surface-induced current with out of plane spin polarization is greater than the normal state current by the factor $\sim T_c E_f^2/\alpha^3$. The maximal amplitude at the surface is solely determined by the structure of the SC gap and formally survives even in the limit $\alpha \to 0$. SC spin currents decay into the bulk on the scale of the coherence length, much slower than in the normal phase. The oscillations in Figs. 4(c) and 4(d) are determined by the spin-orbit strength α and appear due to Faraday-like rotations of the spin coherence functions along quasiparticle trajectories.

In summary, we presented a general analysis of surface bound states and the associated spin currents in noncentrosymmetric superconductors, and applied it to a system with a Rashba-type spin-orbit coupling. We found that the suppression of superconductivity near the surface gives rise to a finite-bias peak in the surface density of states that can be probed by point contact tunneling. We predicted that large in amplitude and slowly decaying spin currents with out of plane spin polarization are generically carried by these surface states. Our prediction opens the route to future investigations and applications of spin transport in systems containing superconductors without center of inversion, and for their use in spin-based devices.

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*Present address: Department of Physics, University of Wisconsin, Madison, WI, USA.

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