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Inverse spin Hall effect in superconductor/normal-metal/superconductor Josephson junctions

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We consider dc supercurrents in SNS junctions. Spin-orbit coupling in combination with Zeeman fields can induce an effective vector potential in the normal conductor. As a consequence, an out-of-plane spin density varying along the transverse direction causes a longitudinal phase difference between the superconducting terminals. The resulting equilibrium phase-coherent supercurrent is analog to the nonequilibrium inverse spin Hall effect in normal conductors. We explicitly compute the effect for the Rashba spin-orbit coupling in a disordered two-dimensional electron gas with an inhomogeneous perpendicular Zeeman field.

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The spin Hall effect (SHE) and inverse SHE (ISHE) are remarkable demonstrations of the influence of the spin-orbit coupling on electron transport. Via this coupling, a longitudinal electric current can induce a perpendicular spin current and vice versa. These effects take place in metals and semiconductors, where the spin-orbit interaction (SOI) arises from impurity scattering[1] or band-structure effects.[2] Utilizing spin injection, SHE, and ISHE, electron spins can be controlled, as recently demonstrated experimentally.[3]

We discuss the intrinsic SHE and ISHE, where the dominant spin-orbit coupling is from the electron band structure. The study of SHE has been focused on normal conductors, e.g., normal metals and semiconductors. Interesting, and rich physics occurs in superconductors where electron transport is dissipationless and the ground state exhibits macroscopic coherence. Some superconductivity induced features of the intrinsic SHE have recently been analyzed in bulk superconductors[4] and superconductor-normal-superconductor (SNS) Josephson junctions.[5] The latter work reveals an equilibrium spin accumulation at lateral sample edges, similar to nonequilibrium spin accumulation in normal conductors, but the spin Hall current vanishes due to time-reversal symmetry in the dc Josephson effect.

We focus on ISHE in Josephson junctions. There are two scenarios depending on how the spin current (density) is created in the normal metal. In a dissipative setup, additional normal/ferromagnetic terminals in the transverse direction inject a nonequilibrium spin current. Subsequently, the ISHE induces an electric potential difference V_{SH} between superconducting terminals, causing Josephson oscillations at frequency $2eV_{SH}/\hbar$. Transport is dissipative due to the spin flow between transverse normal/ferromagnetic terminals. This phenomenon is interesting from an experimental point of view and we will study it quantitatively elsewhere, but we consider here a dissipationless effect.

We present a inverse dissipationless SHE: an out-of-plane equilibrium spin density spatially varying in the transverse direction induces a longitudinal electric supercurrent. Equivalently, it induces a phase shift between two superconducting terminals. In general, since the equilibrium spin density controls ISHE, Zeeman interaction from magnetic or exchange fields manipulates the resulting Josephson supercurrent. As an explicit illustration, we consider the interplay of spin-orbit coupling and Zeeman fields in a disordered two-dimensional electron gas (2DEG), and compute the magnitude of the equilibrium Josephson ISHE.

The interplay of Zeeman field and SOI leading to an effective phase difference between superconducting terminals has recently also been studied in two quite different systems, but neither exhibits the ISHE we discuss: a supercurrent in response to a *spatially homogenous* magnetic field occurs in Josephson tunneling through a one-dimensional (1D) wire[6] and appears in numerical simulations of the superconducting transport through a ballistic point contact[7] in a *spatially homogenous* parallel magnetic field. Note that a normal system analog of the latter phenomenon is the spin-galvanic effect[8] that is different from ISHE. In addition to our main finding of an inverse SHE, we provide an improved understanding of these phenomena by showing how the interplay of Zeeman field and SOI can result in the appearance of an effective electromagnetic vector potential. Such a vector potential, in direct analogy with the Meissner effect, gives rise to a supercurrent.

Let us outline our model. The spin-orbit interaction arises from the band structure, $H_{so}=\boldsymbol{\sigma}\cdot\mathbf{h}_{\boldsymbol{k}}$, where $\boldsymbol{\sigma}=(\sigma_x,\sigma_y,\sigma_z)$ is a vector of Pauli matrices. We assume that the spin-orbit field $\mathbf{h}_{\boldsymbol{k}}$ is given by Rashba SOI where $h_x=\alpha k_y$ and $h_y=-\alpha k_x$ (Ref. 9). Two examples of spin-density manipulations in 2DEG will be considered: (i) a perpendicular to 2DEG Zeeman field spatially varying in the transverse direction y, as shown in Fig. 1 and (ii) a homogeneous Zeeman field directed along the y axis. We will show that setup (i) exhibits an equilibrium inverse SHE. Setup ii) also changes the current-phase relation in SNS contacts. All relevant length scales are assumed larger than the mean-free path $l=v_F\tau$, and we are in the metallic regime $k_F l \gg 1$, where k_F and v_F are the Fermi wave vector and velocity, respectively. These conditions allow a diffusion approximation in the description of electron transport. In this regime, the transport properties are described by a generalized Usadel equation, which we will now derive. The resulting Usadel equation is similar to the one in Ref. 5, but important nontrivial new terms essential for the effects we discuss are added due to the Zeeman interaction $H_Z(\mathbf{r})=\sigma_z H_z(\mathbf{r})+\sigma_y H_y(\mathbf{r})$, where H_z (H_y) are the perpendicular (in-plane) components of the Zeeman field. We start from the anomalous retarded thermal equilibrium Green function $F_{\alpha\beta}(\mathbf{r},\mathbf{k},\omega)$, which is the Fourier transform of

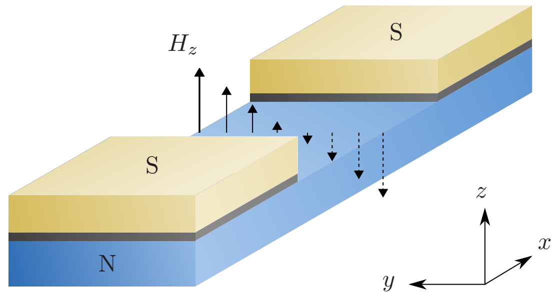

FIG. 1. (Color online) An SNS Josephson junction. Interplay between Rashba spin-orbit interaction and Zeeman splitting in a normal 2D film (n) induces a phase difference between order parameters of two superconducting terminals (s). An inhomogenous Zeeman interaction can be created by, e.g., a ferromagnetic layer on top of the film or magnetic impurities. Another possible configuration (not shown) is a uniform field parallel to the y axis

$$F_{\alpha\beta}=-i\left\langle\left[\psi_\alpha\left(\mathbf{r}+\frac{\boldsymbol{\rho}}{2},t\right),\psi_\beta\left(\mathbf{r}-\frac{\boldsymbol{\rho}}{2},t'\right)\right]_+\right\rangle\theta(t-t') \quad (1)$$

with respect to the relative coordinate $\boldsymbol{\rho}$ and relative time $t-t'$. It is convenient to use a singlet-triplet basis representing the Green's function,

$$F_{\alpha\bar{\beta}}=\frac{1}{\sqrt{2}}(\delta_{\alpha\beta}F_0+\sigma^z_{\alpha\beta}F_s)+\frac{\sigma^+_{\alpha\beta}}{2}F_{+1}+\frac{\sigma^-_{\alpha\beta}}{2}F_{-1}, \quad (2)$$

where $\bar{\beta}$ denotes a spin projection opposite to β, $\sigma^{\pm}_{\alpha\beta}=\sigma^x_{\alpha\beta}\pm i\sigma^y_{\alpha\beta}$. F_s denotes the singlet component. F_0 and $F_{\pm 1}$ are triplet components corresponding to 0 and ± 1 projections of the Cooper pair's total spin on the z axis. Using a standard method starting from Gor'kov equations[10,11] and assuming low SN barrier transmittance, we derive the linearized diffusion equation

$$\sum_m(\delta_{nm}-i\tau K_{nm})F_m=\frac{i}{2\pi N_F}[G^0_{11}\Psi+\Psi G^0_{22}]_n, \quad (3)$$

where subscripts n and m attain the values 0, ± 1, or s, τ is the elastic scattering time and

$$K=2\omega-\mathbf{v}\hat{\mathbf{q}}-2\mathbf{J}\mathbf{h}_{\mathbf{k}}-S-B. \quad (4)$$

Here $\hat{\mathbf{q}}=-i\nabla$, $\mathbf{J}$ is the 3×3 matrix spin 1 operator in the triplet subspace, and operators S and B provide mixing of triplet and singlet components,

$$S_{\pm 1,\mathrm{s}}=-S_{\mathrm{s},\mp 1}=\mp\frac{\hat{\mathbf{q}}}{\sqrt{2}}\frac{\partial h^{\mp}_{\mathbf{k}}}{\partial\mathbf{k}};\ B_{0,\mathrm{s}}=B_{\mathrm{s},0}=2H_z$$

$$B_{\pm 1,\mathrm{s}}=-B_{\mathrm{s},\pm 1}=i\sqrt{2}H_y, \quad (5)$$

where $h^{\pm}=h^x\pm ih^y$. In the right-hand side of Eq. (3) $\Psi=\Sigma_{\mathbf{k}}F$ and the unperturbed retarded Green's functions are

$$G^0_{11/22}=(\omega\mp E_k-\boldsymbol{\sigma}\cdot\mathbf{h}_{\mathbf{k}}-\sigma_zH_z\mp\sigma_yH_y+i\Gamma)^{-1}. \quad (6)$$

The diffusion equation can be derived from Eq. (3) by expanding the operator $(1-i\tau K)^{-1}$ for small τK and averaging over $\mathbf{k}$. The resulting Usadel equation is

$$2i\omega\Psi=\tau\langle(-i\boldsymbol{v}\cdot\nabla+2\mathbf{J}\cdot\mathbf{h}_k)^2\rangle_F\Psi-M\Psi, \quad (7)$$

where the angular brackets denote averaging over the Fermi surface. The matrix M originates from the SOI and the Zeeman interaction expressed via the operators S and B. Its off-diagonal terms describe singlet-triplet transitions. The relevant matrix elements for our further analysis are

$$M_{\mathrm{ss}}=2\tau^3\sum_{\nu=\pm}\nu\langle b^{-\nu}_{\hat{\mathbf{q}}}H_za^{\nu}_{\hat{\mathbf{q}}}+a^{-\nu}_{\hat{\mathbf{q}}}H_zb^{\nu}_{\hat{\mathbf{q}}}\rangle_F,$$

$$M_{\mathrm{s}\pm 1}=\frac{4i\tau^2}{\sqrt{2}}\left\langle 2H_zb^{\pm}_{\hat{\mathbf{q}}}+b^{\pm}_{\hat{\mathbf{q}}}H_z\mp\frac{1}{2}h^2_{\mathbf{k}}a^{\pm}_{\hat{\mathbf{q}}}\right\rangle_F-\sqrt{2}H_y,$$

$$M_{\pm 1\mathrm{s}}=\frac{4i\tau^2}{\sqrt{2}}\left\langle H_zb^{\mp}_{\hat{\mathbf{q}}}+2b^{\mp}_{\hat{\mathbf{q}}}H_z\mp\frac{1}{2}h^2_{\mathbf{k}}a^{\mp}_{\hat{\mathbf{q}}}\right\rangle_F+\sqrt{2}H_y,$$

$$M_{\mathrm{s}0}=M_{0\mathrm{s}}=-2iH_z, \quad (8)$$

where $a^{\pm}_{\hat{\mathbf{q}}}=\hat{q}^i\,\partial h^{\pm}_{\mathbf{k}}/\partial\hat{k}^i$ and $b^{\pm}_{\hat{\mathbf{q}}}=h^{\pm}_{\mathbf{k}}(\mathbf{v}\cdot\hat{\mathbf{q}})$ so that, e.g., the singlet-singlet diagonal element is proportional to ∇_x.

In order to understand some of the underlying physics described by Eq. (7), we will demonstrate that SOI in combination with the Zeeman field gives rise to an effective Meissner effect. Let us first discuss this in the most transparent "local" approximation when the SOI is strong enough/the system long enough, so that the spin-diffusion length $L_{so}=\sqrt{D/\Gamma_{so}}\ll L,\ \sqrt{D/T}$, where L is the length of the junction, $\Gamma_{so}=2\tau\langle h^2\rangle_F$ is the spin-relaxation rate and $D=v_F^2\tau/2$ is the diffusion constant. In this approximation derivatives in triplet parts of Eq. (7) can be disregarded, except in a narrow range $\sim L_{so}$ near the boundaries. H_z is assumed to vary slowly on the L_{so} scale. Expressing the triplet components of Ψ via the singlet Ψ_s and substituting them into the singlet projection of Eq. (7), the latter takes the form

$$2i\omega\Psi_{\mathrm{s}}=-D\nabla^2_x\Psi_{\mathrm{s}}+2iA\nabla_x\Psi_{\mathrm{s}}, \quad (9)$$

where A is a real coefficient obtained from the equation

$$2iA\nabla_x=M_{\mathrm{ss}}+\frac{1}{\Gamma_{so}}\sum_{m=\pm 1}M_{\mathrm{s}m}M_{m\mathrm{s}}. \quad (10)$$

Here we have only included dominant terms proportional to $\alpha^2\partial H_z/\partial y$ and αH_y. Higher order contributions to Eq. (9) proportional to H^2 and α^4 have been disregarded.

The diffusion Eq. (9) demonstrates that cA/e is an effective weak-electromagnetic vector potential. Therefore, similar to the Meissner effect, it will induce a supercurrent. To order A^2 the solution of Eq. (9) is $\Psi_{\mathrm{s}}=\Psi^0_{\mathrm{s}}\exp(ixA/D)$, where Ψ^0_{s} satisfies Eq. (9) with $A=0$. The exponential factor gives rise to an additional phase difference $\theta=LA/D$ between the superconducting terminals, the Josephson current is $j_c\sin(\phi+\theta)$, where ϕ is the initial phase difference between the terminals and j_c is the critical current determined by the function Ψ^0_{s}. The coefficient A is simple for Rashba SOI. For a parallel Zeeman field $A=4\alpha\tau H_y$. For a perpendicular field it vanishes, which is expected since it is similar to the behavior of the spin Hall conductance. Continuing such an analogy, one can expect that $A \neq 0$ for the cubic Dresselhaus[13] SOI.[12]

In order to find a finite ISHE even for the Rashba SOI, we must go beyond the local approximation. In this case, the diffusion Eq. (7) cannot be reduced to simple form (9). We consider superconducting leads with equal real order parameters Δ connected via two SN interfaces with a low transparency t. The barriers are assumed to extend into the 2DEG under the superconducting leads, so that the range of a free-electron motion is between x_L and x_R at the left and right leads, respectively. Depending on contact fabrication, other models can be similarly studied. For example, the electrons in the 2DEG could move freely underneath the contacts, with the barriers present only in z direction, as shown in Fig. 1. The choice of the model is not important for the main qualitative results obtained below.

To the lowest order in the tunneling transparency t, the superconducting current can be expressed[14] as a sum over Matsubara frequencies $\omega = \pi(2n+1)T$,

$$j = \frac{\pi e T}{2R_b^2 N_F}\sum_\omega \frac{\Delta^2}{\Delta^2+\omega^2}\mathrm{Im}\left[\int dy dy' f_{\mathrm{ss}}(\mathbf{r}_L,\mathbf{r}_R')\right], \quad (11)$$

where R_b is the boundary resistance,[15] $\mathbf{r}_{L/R} = (x_{L/R}, y)$ and $f_{ab}(\mathbf{r}_L,\mathbf{r}_R')$, with $a,b=0, \pm 1, s$, is the Green's function of Eq. (7), i.e., a solution of Eq. (7) with a delta source in its right-hand side. The equations for retarded and advanced functions must be properly continued to the upper and lower complex semiplanes of ω, respectively. Treating M in Eq. (7) perturbatively one can express the correction to $f_{\mathrm{ss}}^{(0)}(\mathbf{r}_L,\mathbf{r}_R')$ as

$$\delta f_{\mathrm{ss}}(\mathbf{r}_L,\mathbf{r}_R') = -\int d\mathbf{r} f_{\mathrm{ss}}^{(0)}(\mathbf{r}_L,\mathbf{r})M_{\mathrm{ss}}f_{\mathrm{ss}}^{(0)}(\mathbf{r},\mathbf{r}_R) + \sum_{mm'}\int d\mathbf{r}_1 d\mathbf{r}_2 f_{\mathrm{ss}}^{(0)}$$
$$\times(\mathbf{r}_L,\mathbf{r}_1)M_{sm}f_{mm'}^{(0)}(\mathbf{r}_1,\mathbf{r}_2)M_{m's}f_{\mathrm{ss}}^{(0)}(\mathbf{r}_2,\mathbf{r}_R'), \quad (12)$$

where the unperturbed diffusion propagators $f_{\mathrm{ss}}^{(0)}(\mathbf{r},\mathbf{r}')$ and $f_{mm'}^{(0)}(\mathbf{r},\mathbf{r}')$, with $m,m'=0$ or ± 1, are obtained from Eq. (7) with hard-wall boundary conditions, $\nabla_x f_{\mathrm{ss}}^{(0)} \to 0$ at $x=x_L$ and $x=x_R$, while the triplet components in the case of Rashba SOI satisfy the boundary condition $(iL_{so}\nabla_x + 2J_y)f=0$.[16]

To illustrate the ISHE, we consider the case of Rashba SOI with finite H_z, $H_y=0$, and $\Delta \gg \omega$. The parameter of interest is the effective phase difference between superconducting terminals,

$$\theta_{\mathrm{eff}} = \frac{\sum_\omega \frac{\Delta^2}{\Delta^2+\omega^2}\mathrm{Im}\left[\int dy dy' f_{\mathrm{ss}}(\mathbf{r}_L,\mathbf{r}_R')\right]}{\sum_\omega \frac{\Delta^2}{\Delta^2+\omega^2}\mathrm{Re}\left[\int dy dy' f_{\mathrm{ss}}^0(\mathbf{r}_L,\mathbf{r}_R')\right]}. \quad (13)$$

At $\theta_{\mathrm{eff}} \ll 1$ this parameter allows to express Eq. (11) in the form $j=j_c \sin\theta_{\mathrm{eff}}$. From Eq. (7), $\theta_{\mathrm{eff}} = C\rho$, where $C = 2l\overline{\nabla_y H_z}\tau/k_F L$ and $\overline{\nabla_y H_z}$ denotes the average value of the magnetic field gradient in the contact range. ρ is shown in Fig. 2 as a function of ratio of the spin-relaxation rate $\Gamma_{so} = 2\tau\alpha^2 k_F^2$ versus the Thouless energy $E_T = D/L^2$. For large SOI the "local" approximation is obtained by using the approximate form of $f_{\pm1\pm1} = 2f_{00} = -\delta(\mathbf{r}_1-\mathbf{r}_2)/\Gamma_{so}$ in the second term of Eq. (12). In this case both terms in Eq. (12) are proportional to α^2 and precisely cancel each other, as in the r.h.s. of Eq. (10). Beyond this leading "local" approximation there are terms increasing slower than α^2. They contribute to Fig. 2.

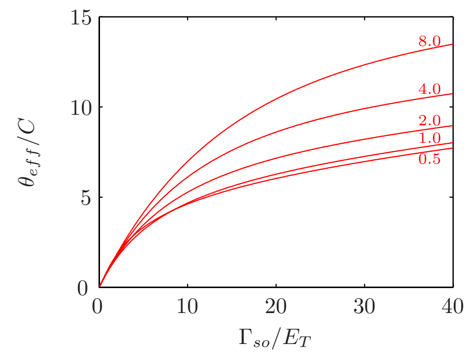

FIG. 2. (Color online) The phase difference versus a ratio of the spin-relaxation rate and the Thouless energy, at $0.5 < k_B T/E_T < 8$. The parameter C is described in the text.

Larger Zeeman fields cannot be treated perturbatively. A strong depairing effect takes place when the characteristic length $L_Z = \sqrt{D/2H}$ is small, $L_Z \ll \min(L, L_{so})$. Then, for $H = H_z$, both Ψ_s and Ψ_0 decay exponentially near contacts with superconducting terminals and the latter become effectively disconnected. On the other hand, as it follows from Eq. (7), $\Psi_{\pm1}$ components are not subject to the depairing effect and can propagate at the relatively large distance $\sim L_{so}$. Such a long-range triplet effect has been studied in SFS junctions, where a link between triplet and singlet Cooper pairs has been induced by an inhomogeneous (rotating) magnetization (see Ref. 11 and references therein). In our case, a coupling of $\Psi_{\pm1}$ to Ψ_0 and Ψ_s can be provided by SOI through the matrix elements $M_{\pm1s}$ and the spin precession operator $R_{\pm1,0} = -i4\tau\langle(\mathbf{J}_{\pm1,0}\cdot\mathbf{h}_k)(\boldsymbol{v}\cdot\nabla)\rangle_F$ originating from the first term in the right-hand side of Eq. (7). Indeed, assuming that $H_z \gg \Gamma_{so}$ and E_T, it is easy to show that modified Eq. (12) is represented by its second term, where the integrand has the form

$$f_{s0}^{(0)}(x_L,x)R_{0m}f_{mm}^{(0)}(x,x')M_{ms}f_{\mathrm{ss}}^{(0)}(x',x_R). \quad (14)$$

The unperturbed functions $f_{s0}^{(0)}$ and $f_{\mathrm{ss}}^{(0)}$ are obtained from s and 0 projections of Eq. (7), where the precession term and all M_{ij}, except M_{s0} and M_{0s} are ignored. The physics of the process described by Eq. (14) is clear: the magnetic field mixes 0-triplet and singlet components of the pairing function within the short range near the left boundary. Further, due to the spin precession in the SOI field the 0-triplet transforms to ±1 triplet components. The latter propagate to the right contact where they convert to the singlet through $M_{\pm1s}$. Integrating Eq. (14) over x and x' gives a power-law dependence of $\mathrm{Im}[\delta f_{\mathrm{ss}}(x_L,x_R)]$ on the magnetic field,

$$\mathrm{Im}[\delta f_{\mathrm{ss}}(x_L,x_R)] \propto (|H_z(y_1)|^{-3/2} - |H_z(y_2)|^{-3/2}), \quad (15)$$

where y_1 and y_2 are y coordinates of the junction edges. Accordingly, at $L_{so} \ll L$ and $|y_1-y_2| \sim L$ an order-of-magnitude evaluation of θ_{eff} can be written as

$$\theta_{\text{eff}} \simeq \frac{l^3}{2L_{so}^2 L} \frac{\rho}{k_F l}\left(\left| \frac{H_c}{H_z(y_1)} \right|^{3/2} - \left| \frac{H_c}{H_z(y_2)} \right|^{3/2} \right), \quad (16)$$

where $2H_c = D/L_{so}^2 = \Gamma_{so}$ and ρ is shown at Fig. 2.

In contrast to a perpendicular Zeeman field, in a parallel field ± 1 triplets exponentially decay near boundaries, as can be seen from Eqs. (7) and (8). So they cannot provide a long-range link between superconducting terminals.

In conclusion, an analog to the ISHE exists in dc Josephson SNS junctions. Unlike the normal ISHE, the supercurrent through the SNS contact can be induced by a static Zeeman interaction by magnetic or exchange fields oriented normal to the 2DEG and varying in the direction transverse to the electric current. A destructive depairing effect of the strong Zeeman field is diminished by Rashba SOI leading to a power-low dependence on this field. We show that a supercurrent through the junction can also be induced by a uniform parallel Zeeman field, corroborating thus the numerical analysis of Ref. 7. On the other hand, the depairing effect of such a field was found to be strong (exponential). In both cases an appearance of the supercurrent can be explained in terms of the Meissner effect produced by an effective vector potential, which is a combined effect of the Zeeman field and Rashba spin-orbit interaction.

We considered the diffusive transport regime which is relevant in low mobility metals and (magnetic) semiconductors. Furthermore, the diffusive regime, allows an elucidation of the main physics and parameters governing this phenomena. We expect a strong Josephson ISHE in ballistic junctions containing a metallic normal layer with a strong Rashba interaction, for example in Bi films on some substrates.[17] Ballistic quantum wells of narrow gap semiconductors are also expected to exhibit an increased Josephson ISHE.

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