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Spin-resolved Andreev reflection in ferromagnet-superconductor junctions with Zeeman splitting

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Abstract. – Andreev reflection in ferromagnet-superconductor junctions is derived in a regime in which Zeeman splitting dominates the response of the superconductor to an applied magnetic field. Spin-up and spin-down Andreev reflections are shown to be resolved as the voltage is increased. In the metallic limit, the transition from Andreev to quasiparticle transport in the spin-up channels has a non-trivial behavior when spin polarization increases. The conductance is asymmetric in a voltage reversal, which can be used as a new probe of spin polarization.

The interplay between Andreev reflection and spin polarization has generated recently an important interest, both theoretical [1–8] and experimental [9–15]. The subgap conductance in normal-metal–superconductor (NS) junctions originates from Andreev reflection [16]: a spin-σ electron incoming from the N side is reflected as a hole in the spin-$(-\sigma)$ band while a spin-zero Cooper pair is transferred into the superconductor. Since the incoming electron and outgoing hole belong to opposite spin bands, Andreev reflection couples to a Fermi surface polarization in the N side of the junction. de Jong and Beenakker [1] showed theoretically that increasing the Fermi surface polarization in ferromagnet-superconductor (FS) junctions suppresses Andreev reflection because Andreev reflection is limited by the minority-spin channels. Their prediction was verified experimentally by Soulen *et al.* [14] and Upadhyay *et al.* [15], who used this effect to measure the Fermi surface polarization. On the other hand, Tedrow and Meservey [17] demonstrated that under specific conditions, a magnetic field can be used to tune a Zeeman splitting of the quasiparticle excitations in a superconductor [17], and used it to perform a spin-resolved tunnel spectroscopy in FS junctions [17]. I show in this letter that Zeeman splitting can be used to resolve the spin-up and spin-down Andreev reflections, with a different threshold voltage $eV_\pm = \Delta \mp \mu_\mathrm{B} H$ for the transition from Andreev to quasiparticle transport in a magnetic field H. In NS junctions with Zeeman splitting, the spin-up and spin-down differential conductances have the same behavior at the Andreev reflection threshold voltages $V_\pm$. In FS junctions with Zeeman splitting, a non-trivial behavior at the spin-up threshold voltage V_+ is predicted. In addition, the conductance is asymmetric in a voltage reversal, which allows a possible new determination of the Fermi surface spin polarization based on spin-resolved Andreev reflection.

Our modeling is intented to describe a point contact experiment in which the transverse dimension is smaller than the scattering length in the ferromagnetic metal. We neglect disorder in the superconductor, as well as the proximity effect in the N or F sides of the junction [12, 18–26]. Such a treatment captures the interplay between spin polarization and Andreev reflection in a multichannel point contact, which appears to be the relevant physics in recent experiments [14, 15]. Moreover, Soulen *et al.* [14] succeeded in realizing a high-transparency point contact between a superconductor and a ferromagnet by a mechanical adjustment. A similar technique may be used to experiment the situation we consider in this letter.

Let us first consider a NS point contact and derive the physics associated with the interplay between Andreev reflection and Zeeman splitting in the superconductor. The superconductor is assumed to have a thin film geometry with the magnetic field applied parallel to the film. We assume a small orbital depairing parameter while the critical field for destroying superconductivity is set by Pauli paramagnetism [27], with large values of $H_{c2\|} \sim 5$ T for Al thin films [17]. The spin-orbit scattering length is supposed to be small compared to the superconductor coherence length ξ, as is the case for light elements such as Al [17]. This insures that electrons in the superconductor have a well-defined spin σ at length ξ, and therefore a well-defined Zeeman energy $-\mu_{\rm B} H\sigma$ [17, 28]. The coherence factors of spin-σ electrons (u_σ) and holes in the spin-($-\sigma$) band ($v_{-\sigma}$) with an energy ϵ are

$$u_\sigma^2 = 1 - v_{-\sigma}^2 = \frac{1}{2}\left(1 + \frac{\sqrt{(\epsilon+\sigma\mu_{\rm B}H)^2 - |\Delta|^2}}{\epsilon+\sigma\mu_{\rm B}H}\right), \tag{1}$$

with therefore a coupling between Andreev reflection and Zeeman splitting. A step function variation of the superconducting gap at the interface is assumed: $\Delta(x) = \Delta\theta(x)$. We consider a δ-function elastic interface scattering potential $V(x) = H_0\delta(x)$, interpolating between a metallic contact if $H_0 = 0$ and a tunnel junction if $H_0 = \infty$ [29]. The interface barrier is normalized with respect to the Fermi velocity: $Z = mH_0/(\hbar\sqrt{2m\mu})$, with $\mu = \hbar^2 k_{\rm F}^2/2m$ the chemical potential [29]. The energy dependence of the transmitted quasiparticle wave vectors is irrelevant to the present calculation [30]. We assume the normal metal/ferromagnet and the superconductor to have identical band widths. The presence of a different band width would renormalize the conductance [5] while keeping the possibility of a spin-resolved Andreev reflection unchanged.

Since multichannel effects in the NS junction model do not play the same crucial role as in FS junctions, we first focus on the single-channel NS junction model. The coherence factors equation (1) leads to the Andreev reflection transition probability of electrons with a spin-σ and holes in the spin-($-\sigma$) band with an energy ϵ:

$$A^{e\uparrow}(\epsilon) = A^{h\downarrow}(\epsilon) = A_{\rm BTK}(\epsilon+\mu_{\rm B}H), \quad \text{and} \quad A^{e\downarrow}(\epsilon) = A^{h\uparrow}(\epsilon) = A_{\rm BTK}(\epsilon-\mu_{\rm B}H), \tag{2}$$

with $A_{\rm BTK}(\epsilon)$ the Blonder, Thinkham and Klapwijk (BTK) Andreev reflection coefficient [29]. Equations (2) are valid also if $\epsilon < 0$, in which case transmission of quasiparticles on negative energy branches should be considered. The zero-temperature differential conductance of spin-σ carriers in the presence of Zeeman splitting is $G^\sigma(eV, H) = G_{\rm BTK}(eV + \sigma\mu_{\rm B}H)$, with $G_{\rm BTK}(eV)$ the BTK differential conductance in the absence of Zeeman splitting [29]. $G^\sigma(eV, H)$ shows a Zeeman splitting for an arbitrary interface scattering in the sense that the magnetic field enters the conductance via the combination $eV + \sigma\mu_{\rm B}H$ only. The tunnel

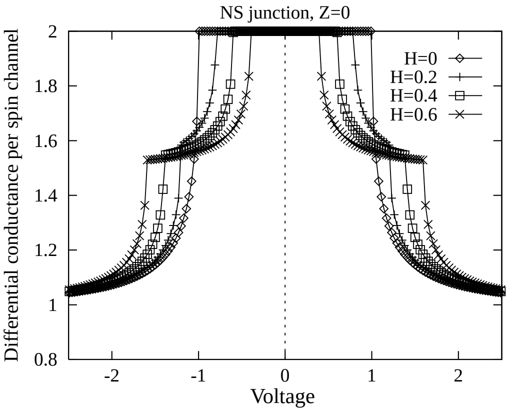

Fig. 1 – Differential conductance of the NS junction in the metallic limit $Z = 0$, with a Zeeman splitting $\mu_B H = 0$ ($\Diamond$), 0.2 (+), 0.4 ($\Box$), 0.6 ($\times$) in units of the superconducting gap Δ. The conductance is normalized to the number of spin channels. The voltage is in units of the superconducting gap Δ. A plateau of $3e^2/(2h)$ per spin channel develops in the conductance when H increases. We have considered in this figure a multichannel model generalizing eq. (2), with $N^\uparrow = N^\downarrow = 141$ channels and $\mu = 10^4$. The behavior of the multichannel NS model is identical to the single-channel NS model.

spectrum in the limit $Z \gg 1$ reproduces the Zeeman splitted density of states of the superconductor $\rho_\sigma(\epsilon) = \rho_{\rm BCS}(\epsilon + \sigma\mu_B H)$, with $\rho_{\rm BCS}$ the single-spin BCS density of states [17,28,29]. In the metallic limit $Z = 0$ and below the spin-up threshold voltage $eV_+ = \Delta - \mu_B H$, spin-up and spin-down transport originates from Andreev reflection, with a conductance equal to $2e^2/h$ per spin channel (see fig. 1). Spin-up transport transits from Andreev reflection to quasiparticle transport at the spin-up threshold voltage eV_+, smaller than the spin-down threshold voltage $eV_- = \Delta + \mu_B H$. In between V_+ and V_- a plateau of $3e^2/(2h)$ per spin channel develops in the conductance when H increases, corresponding to an Andreev reflection transport of spin-down carriers and a quasiparticle transport of spin-up carriers. The strongest field $\mu_B H = 0.6\Delta$ in fig. 1 is close to the paramagnetic breakdown. Figure 1 tends to indicate that a well-defined plateau of $3e^2/(2h)$ may not be observed experimentally. Nevertheless, two distinct features at $eV_\pm$ are already present even for moderate magnetic fields (for instance, $H = 0.2\Delta$ in fig. 1).

We now discuss the effect of a spin polarization in the normal metal, in which case multichannel effects play a relevant role. We show a non-trivial transition from Andreev to quasiparticle transport at the spin-up threshold voltage V_+, as well as a conductance asymmetric in a voltage reversal. We denote by n and n' the quantum numbers associated to a quantized transverse motion in a clean FS point contact of cross-sectional area a^2. We assume a Stoner ferromagnet with an exchange field $h_{\rm ex}(x) = h_{\rm ex}\theta(-x)$. The channel with transverse quantum numbers (n, n') in the spin-σ band has a dispersion $E^\sigma_{n,n'}(k^\sigma) = \hbar^2(k^\sigma)^2/(2m) - \sigma h_{\rm ex} + \kappa(n^2 + n'^2)$, with the energy $\kappa = (\hbar^2/2m)(\pi/a)^2$ inverse proportional to the junction area, and related to the number of spin-σ channels according to $N^\sigma = \pi(\mu + \sigma h_{\rm ex})/(4\kappa)$ [15]. The associated barrier parameter $Z^{F,\sigma}_{n,n'}$ of spin-σ electrons in the channel (n, n') is $Z^{F,\sigma}_{n,n'} = \left(1 + \sigma\frac{h_{\rm ex}}{\mu} - \frac{\kappa}{\mu}(n^2 + n'^2)\right)^{-1/2} Z$, with $Z = mH_0/(\hbar\sqrt{2m\mu})$. The transverse dimensions of the S side of the junction are assumed to be identical to the ones of the N side and the gap, the interface scattering and the exchange field are constant in the transverse direction, with therefore a conservation of the transverse quantum numbers across

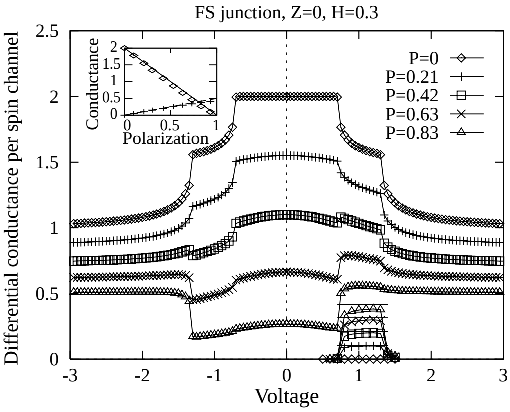

Fig. 2 – Differential conductance of the FS junction in the metallic limit $Z = 0$. The conductance is normalized to the number of spin channels. The voltage is in units of the superconducting gap Δ. The chemical potential is $\mu = 10^4$, and the Zeeman splitting is $\mu_B H = 0.3$ (in units of Δ). In the absence of a spin polarization, we have $N^\uparrow = N^\downarrow = 141$ channels. The Fermi surface polarizations $P = (N^\uparrow - N^\downarrow)/(N^\uparrow + N^\downarrow)$ are $P = 0$ ($\Diamond$), $P = 0.21$ ($+$), $P = 0.42$ ($\Box$), $P = 0.63$ ($\times$), and $P = 0.83$ ($\triangle$). The conductance is asymmetric in a voltage reversal, and has a non-trivial behavior at the spin-up threshold voltage $eV_+ = \Delta - \mu_B H$. The bottom curves show the asymmetry of the conductance spectrum $G_A(V) = G(V) - G(-V)$ in the range $V_+ < V < V_-$ for the same values of the spin polarization. A plateau at $G_A(V) \simeq P/2$ is present (solid lines). The insert shows the zero-voltage conductance ($\Diamond$) and the asymmetry ($+$) *vs.* spin polarization P, compared to their values $2(1-P)$ and $P/2$ obtained from a channel counting argument.

the interface [31]. The pairing Hamiltonian in the S side with a cross-sectional area a^2 is

$$H^S = \sum_{n,n',k,\sigma} \left(\frac{\hbar^2 k^2}{2m} + \kappa(n^2 + n'^2) \right) c^+_{n,n',k,\sigma} c_{n,n',k,\sigma} + \sum_{n,n',k} \left(\Delta c^+_{n,n',k,\uparrow} c^+_{n,n',-k,\downarrow} + \text{h.c.} \right),$$

with an associated barrier parameter $Z^{\rm S}_{n,n'} = \left(1 - \frac{\kappa}{\mu}(n^2+n'^2)\right)^{-1/2} Z$, different from $Z^{{\rm F},\sigma}_{n,n'}$ because of the exchange field generating different Fermi wave vectors in the ferromagnet and the superconductor. The channels with a spin-up Fermi surface only have a real positive $Z^{{\rm F},\uparrow}_{n,n'}$ and a pure imaginary $Z^{{\rm F},\downarrow}_{n,n'}$. Physically, a spin-up electron incoming from the N side below the superconducting gap in such a channel is Andreev-reflected into an evanescent hole state in the spin-down band, with a pure imaginary wave vector $k^\downarrow$. The hole propagates in the ferromagnet over the length scale $1/\text{Im}(k^\downarrow)$ before it is backscattered onto the interface and Andreev-reflected as a spin-up electron, therefore not carrying current, as proposed by de Jong and Beenakker [1]. Incorporating this process under the form of a pure imaginary interface scattering allows to calculate transport above the superconducting gap. The matching of the wave functions between the F and S sides is solved similarly to refs. [29], including the coherence factors in eq. (1), and the barrier parameters $Z^{{\rm F},\sigma}_{n,n'}$ and $Z^{\rm S}_{n,n'}$ [32]. The resulting differential conductance spectra are shown in fig. 2 in the metallic limit $Z = 0$.

At low voltage, the conductance shows a reduction of Andreev reflection by spin polarization [1]. The large voltage limiting value of the quasiparticle conductance per spin channel decreases from the Landauer value e^2/h in the absence of spin polarization to $e^2/(2h)$ with a full polarization, because only the ferromagnet channels with a corresponding channel in

TABLE I – *Number of two reflection and quasiparticle transport channels. "Spin-σ AR" stands for spin-σ Andreev reflection. "Spin-σ QP" stands for spin-σ quasiparticle transport. The conductance (in units of e^2/h) is normalized to the number of spin channels.*

Voltage	$-V_- < V < -V_+$	$-V_+ < V < V_+$	$V_+ < V < V_-$	$V_- < V$ $V < -V_-$
Spin-up AR	$N^\downarrow$ channels	$N^\downarrow$ channels	–	–
Spin-down AR	–	$N^\downarrow$ channels	$N^\downarrow$ channels	–
Spin-up QP	–	–	N^{S} channels	N^{S} channels
Spin-down QP	$N^\downarrow$ channels	–	–	$N^\downarrow$ channels
Conductance	$3/2 - 3h_{\mathrm{ex}}/(2\mu)$	$2 - 2h_{\mathrm{ex}}/\mu$	$3/2 - h_{\mathrm{ex}}/(2\mu)$	$1 - h_{\mathrm{ex}}/(2\mu)$

the superconductor contribute to the quasiparticle conductance. The number of spin-down quasiparticle transport channels is $N^\downarrow$, while spin-up quasiparticle transport is limited by the number of superconducting channels $N^{\mathrm{S}} = \pi\mu/4\kappa$. The total number of quasiparticle transport channels is therefore $\pi(2\mu - h_{\mathrm{ex}})/4\kappa$, reduced by a factor of two when the exchange field h_{ex} increases from zero to μ.

Now the behavior of the differential conductance at the spin-up threshold voltage eV_+ differs qualitatively in the weak and strong polarization regimes: the conductance decreases with voltage at eV_+ if spin polarization is weak while it increases if spin polarization is strong (see fig. 2). With a weak polarization, most of the spin-up channels are Andreev-reflected and the decrease in conductance at eV_+ can be understood qualitatively on the basis of the transition from Andreev to quasiparticle transport in the single-channel BTK model [29]. If spin polarization is strong, a fraction $1 - (N^\downarrow/N^\uparrow)$ of the spin-up channels are not Andreev-reflected if $V < V_+$. These channels however contribute to the quasiparticle current if $V > V_+$, with a spin-up quasiparticle conductance $\simeq (e^2/h)N^{\mathrm{S}}$, larger than the Andreev conductance

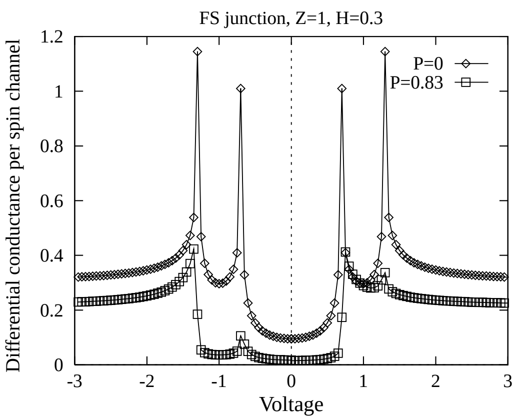

Fig. 3 – Differential conductance of the FS junction with $Z = 1$. The conductance is normalized to the number of spin channels. The voltage is in units of the superconducting gap Δ. The chemical potential is $\mu = 10^4$, and the Zeeman splitting is $\mu_{\mathrm{B}}H = 0.3$ (in units of Δ). In the absence of spin polarization, we have $N^\uparrow = N^\downarrow = 141$. The Fermi surface polarizations $P = (N^\uparrow - N^\downarrow)/(N^\uparrow + N^\downarrow)$ are $P = 0$ ($\Diamond$), and $P = 0.83$ ($\Box$). Spin polarization results simultaneously in a suppression of Andreev reflection and spin-polarized tunneling.

$\simeq (2e^2/h)N^{\downarrow}$ if $h_{\rm ex} > \mu/2$. The contribution of the different spin channels to the conductance is given in table I. The approximate conductances in table I reproduce the correct features of the conductance spectrum obtained in the full calculation in fig. 2. In particular, the asymmetry in the conductance $G_{\rm A}(V) = G(V) - G(-V)$ is equal to $h_{\rm ex}/(2\mu) = P/2$ in the voltage range $V_+ < V < V_-$. As shown in fig. 2, the value $P/2$ of the asymmetry compares well with the full calculation. The asymmetry has the same origin in the low- and high-transparency cases (time-reversal symmetry breaking of quasiparticles in the superconductor due to Zeeman splitting).

Finally, we have shown in fig. 3 the behavior of the FS junction model with a moderate transparency ($Z = 1$). In this parameter range, and in the absence of spin polarization, two tunnel-like peaks coexist with a finite low-voltage conductance originating from Andreev reflection. Increasing spin polarization results in a suppression of Andreev reflection by spin polarization *and* spin-polarized tunneling (a spin-up peak at eV_+ with a stronger weight than the spin-down peak at eV_-). These two phenomena may therefore be observed simultaneously.

To conclude, we have shown that Zeeman splitting can be used to resolve the spin-up and spin-down Andreev reflections in NS and FS junctions. In metallic FS junctions, the spin-up quasiparticle current is larger than the spin-up Andreev reflection current if spin polarization is large. The different behavior in the spin-up and spin-down channels generates a conductance spectrum asymmetric in a voltage reversal, which provides a new possibility to probe the Fermi surface spin polarization. We have also shown that the point contacts with an intermediate interfacial scattering $Z \sim 1$ show simultaneously a reduction of Andreev reflection by spin polarization and spin-polarized tunneling. Finally, it appears that spin-resolved Andreev reflection in ferromagnet-superconductor junctions allows two simultaneous determinations of the Fermi surface polarization of the ferromagnet: one determination from the asymmetry of the conductance and another determination from the zero bias conductance. The model analyzed in this letter leads to a simple relation between these two determinations (see the insert in fig. 2). One may ask whether a similar relation would also hold in an experiment.

* * *

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[30] We calculated the Andreev reflection coefficient of NS junctions with Zeeman splitting with the full dependence of the quasiparticle wave vectors on H. We neglect the resulting correction of order $\mu_{\mathrm{B}}H/\mu \ll 1$ to the Andreev reflection coefficient.
[31] We expect the same qualitative physics as in a contact with different transverse dimensions of the N and S sides. The asymmetric contact may be solved via a recursion transfer matrix method; see Kobayashi N., Takahashi S. and Maekawa S., Report No. cond-mat/9906406 and references therein.
[32] In the energy range $-\Delta + \mu_{\mathrm{B}}H < \epsilon < \Delta - \mu_{\mathrm{B}}H$, the backscattering transition probability is unity, with a real $Z^{\mathrm{F},\uparrow}_{n,n'}$ and a pure imaginary $Z^{\mathrm{F},\downarrow}_{n,n'}$ while Z is arbitrary. This shows the total backscattering of spin-up electrons incoming below the superconducting gap in a channel with no spin-down Fermi surface.