Ferromagnetic resonance with a magnetic Josephson junction

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Received 4 October 2010, in final form 7 December 2010
Published 19 January 2011
Online at stacks.iop.org/SUST/24/024020

Abstract

We show experimentally and theoretically that there is a coupling via the Aharonov–Bohm phase between the order parameter of a ferromagnet and a singlet, s-wave, Josephson super-current. We have investigated the possibility of measuring the dispersion of such spin-waves by varying the magnetic field applied in the plane of the junction and demonstrated the electromagnetic nature of the coupling by the observation of magnetic resonance side-bands to microwave induced Shapiro steps.

(Some figures in this article are in colour only in the electronic version)

Rotation symmetry associated with the O(3) orthogonal group forbids a coupling between a scalar s-wave order parameter and the vector order parameter \( \vec{M} \) of the ferromagnet. However, a Josephson junction [1] defines a plane, the O(3) symmetry is broken, and such a coupling is possible. Here we describe a part of the rich spectroscopic magnetic resonance possibilities that this observation implies. It is possible to perform a ‘photon free’ ferromagnetic resonance (FMR) experiment [2] on about \( 10^7 \) Ni atoms, something unfeasible with standard FMR techniques.

Interactions in nature reflect certain gauge groups and the associated phases which generate a vector potential \( \vec{A} \), called the Berry connection [3]. Interactions via electromagnetic fields generated by the conserved electrical currents reflect the \( U(1) \) gauge group and the Aharonov–Bohm (AB) phase [4]. Associated with the angular moment is the gauge symmetry \( SU(2) \) and the familiar Lie algebra, i.e. the spin commutation rules. The AB phase is replaced by the spin Berry phase. It is often imagined that magnetic moments might interact with the Josephson current by direct spin-flips, which would involve the spin Berry phase [5]. In this paper it will be shown that quantitatively the interaction between the Josephson current and the order parameter in superconductor/ferromagnet/superconductor (SFS) junctions can be explained in terms of the AB phase and regular electrodynamics. It will be shown that such an experiment measures fairly directly the magnet correlation function.

We fabricated Nb/Pd\(_{0.9}\)Ni\(_{0.1}\)/Nb Josephson junctions by \( \textit{in situ} \) angle evaporation through a resist mask and subsequent lift-off. The mask is defined by e-lithography on a polyether sulfone—PES (500 nm)/Si\(_3\)N\(_4\) (60 nm)/polymethyl methacrylate—PMMA (350 nm) trilayer [6], and etched in a reactive ion etching (RIE) chamber. The Si\(_3\)N\(_4\) is etched for 1 min 30 s in SF\(_6\) and the PES in oxygen plasma for 10 min, giving an undercut of 500 nm. The mask fabrication process is schematically presented in figure 1(a). The Si\(_3\)N\(_4\) suspended bridge allows shadow evaporation. A scanning electron microscope (SEM) picture of the mask including the suspended bridge is reported in figure 1(c). The first Nb layer is evaporated at \( -45° \) while the PdNi is evaporated at 45° and the second Nb layer at 47°. The overlap in the \( y \)-direction defines the junction area. The shadow evaporation is illustrated in the lower drawing array of figure 1(a). Eight junctions are evaporated on the sample chip. An SEM image of one of the junctions after lift-off is shown in figure 1(d). The junction area is \( 0.5 \times 0.5 \mu m^2 \). The electron-gun evaporation is carried out in ultra-high vacuum (UHV) with a base pressure lower than \( 10^{-9} \) mbar. The two Nb superconducting films are 50 nm thick and the ferromagnetic layer of Pd\(_{0.9}\)Ni\(_{0.1}\) is 20 nm thick. The Ni concentration is measured by Rutherford backscattering...
junction to junction is about on the same wafer, the dispersion of the critical current from temperature. We have measured four ferromagnetic junctions lift-off.

Figure 1. Description of the sample: (a) the fabrication procedure; (b) the principle of the experiment; (c) SEM picture of the SiN/NabPES mask before angle evaporation; (d) SEM picture of a junction after lift-off.

(RBS) on a test sample. The magnetization loops obtained by superconducting quantum interference device (SQUID) magnetometry with in-plane and out-of-plane magnetic field show a predominant perpendicular anisotropy. Finally, a schematic view of the junction including the principle of the experiment is to be found in figure 1(b).

Typical current–voltage (IV) characteristics are shown in figure 2 as a function of temperature. Well below the critical temperature, the critical current $I_c \approx 7 \mu A$ and normal resistance $R_n = 1 \Omega$ give a Josephson coupling of $7 \mu V$, consistent with early studies on highly underdamped PdNi-based Josephson junctions [7]. The Josephson penetration depth is about 10 $\mu m$. The IV characteristics are not hysteretic, confirming overdamped phase dynamics, and are well described by the resistively shunted junction (RSJ) model [1]. The critical current versus temperature (see inset of figure 2) shows the quasi-linear dependence expected in SFS junctions for $d_F > \xi_F$ and $E_{ex} \gg \Delta > k_B T$ [8]. Where $d_F$ and $\xi_F$ are the thickness and coherence length in the ferromagnetic layer, respectively. The exchange energy, $E_{ex}$, as estimated from the decay of the superconducting order parameter, is $35$ meV [7] and hence is indeed much larger than the superconducting energy gap $\Delta = 1.5$ meV. This quasi-linear temperature dependence has been observed in Nb/CuNi/Nb Josephson junctions with comparable $d_F$ and $E_{ex}$ values [9]. This linear dependence has been observed previously in highly underdamped junctions [7]. The junction critical temperature is about 7.0 K while the critical temperature of the Nb leads is 7.6 K and their critical current is over 500 $\mu A$ at low temperature. We have measured four ferromagnetic junctions on the same wafer, the dispersion of the critical current from junction to junction is about $\pm 1 \mu A$, $\Delta R_n$ is 0.15 $\Omega$ while the $I_c R_n$ varies by less than 3% from junction to junction. We have also fabricated non-magnetic junctions by the same process but replacing the PdNi with a thicker 70 nm Pd layer. These junctions have a much larger critical current of about 44 $\mu A$.

For a square junction of side $L$, the total super-current is given by the integral [1]

$$I_s = J_c \int_{-L/2}^{L/2} dx \int_{-L/2}^{L/2} dy \sin \gamma(x, y; t)$$

(1)

involving [10] the gauge invariant phase:

$$\gamma(x, y; t) = \Delta \varphi(x, y; t) - \frac{2e}{\hbar} \int A \cdot d\mathbf{r},$$

(2)

where $\Delta \varphi(x, y; t)$ is the local difference in the phase of the order parameter and where the last term is the AB phase, involving the vector potential $\mathbf{A}$. Decomposing $\mathbf{A} = A^0 + A^1$, where $A^0$ and $A^1$ reflect the static and dynamic magnetic fields, in the usual way [1],

$$\gamma(x, y; t) = \phi_0 + k x + \omega t - \frac{2e}{\hbar} \int A^1 \cdot d\mathbf{r},$$

where $\phi_0$ is an arbitrary phase, $\omega t = (2e/\hbar) V_0$ is the Josephson frequency, and the wavevector $k = (4e\omega/\hbar) \mu_0 H + (4e\omega/\hbar) \mu_0 M_0$. Here $M_0$ is the $y$-component of the static magnetization $\mathbf{M}_0$ and the applied field $\mathbf{H} = H \hat{y}$. These terms involve, respectively, the actual 2x and magnetic 2$d = 2(a + \lambda)$ widths, where $\lambda$ is the London penetration depth, since $\mathbf{B}$ only differs from $\mu_0 H$ in the magnetic layer. These equations describe both the statics and the dynamics of our junctions. Experimentally, with $H = 0$, $M_0 = M_0 k$, so, reflected in the above, is the hysteretic component $M_0$ induced by $\mathbf{H}$.

Figure 2. Current–voltage characteristics with increasing temperature. The curves have been shifted vertically for clarity. They do not show any hysteresis as expected when the phase dynamics is strongly damped. The critical current as a function of the temperature is shown in the inset, the dashed line is a linear fit.
In linear response, the dynamics fields contained in $\frac{2e}{\hbar} \int A^i \cdot \mathrm{d}r$ are considered as a perturbation. Expanding, the dc signal is

$$I_1 = -\frac{2e}{\hbar \omega_0 T} \int_0^T \mathrm{d}t \int_0^{2a} \mathrm{d}z \int_0^{L/2} \mathrm{d}x \int_0^{L/2} \mathrm{d}y \frac{\partial J_x}{\partial t} \cdot A^i$$

where $\int \mathrm{d}v = \int_0^{L/2} \mathrm{d}x \int_0^{L/2} \mathrm{d}y \int_0^{2a} \mathrm{d}z$ is the integral over the barrier volume, which extends from $z = 0$ to $2a$, and where $J_x = J_z \sin(\phi_0 + kx + \omega t)\hat{z}$ is the appropriate Josephson current density. There is a time average over a single period $T$. This key result is gauge invariant.

The approach is similar to that used for junctions with magnetic impurities in a normal metal barrier [11]. It is necessary to determine the appropriate boundary conditions for the solutions of Maxwell’s equations. For reasons of transparency, it is not at all useful to solve the very difficult problem in which the solution within the barrier is matched to that in the exterior region to the junction. Within the junction we ignore the displacement and transport currents since the wavelength of light $\lambda$ and the skin depth $\delta$ are both larger than the dimensions of the junction at the Josephson frequency $\omega_0$ relevant for the FMR. We observe that the impedance of the barrier volume, which extends from $z = 0$ to $2a$, and within the junction, is essentially no radiation from the junction. The displacement current, and evidently the transport current, can therefore also be ignored in the exterior region. In order to evaluate $H$ it is therefore only necessary to integrate Ampère’s circulation law

$$\nabla \times H = J_z.$$  \hspace{1cm} (4)

However then, in order to evaluate $A$ and hence equation (3), it is also required that $\nabla \times A = B$ be integrated with $B = \mu_0(H + M)$. Even with the present simplifications, this is an involved calculation. It is useful to make some formal manipulations in order to avoid this double integration. Performing an integration by parts with time

$$I_1 = -\frac{2e}{\hbar \omega_0 T} \int_0^T \mathrm{d}t \int_0^{2a} \mathrm{d}z \int_0^{L/2} \mathrm{d}x \int_0^{L/2} \mathrm{d}y \frac{\partial A^i}{\partial t} \cdot \nabla \times H$$ \hspace{1cm} (5)

Using the fact that the Poynting vector and hence $H \times \partial A/\partial t = 0$ on the surface, this is integrated again by parts using $\nabla \cdot (H \times (\partial A/\partial t)) = (\partial A^i/\partial t) \cdot \nabla \times H = H \cdot \nabla \times (\partial A/\partial t)$ to give

$$I_1 = -\frac{2e}{\hbar \omega_0 T} \int_0^T \mathrm{d}t \int_0^{2a} \mathrm{d}z \int_0^{L/2} \mathrm{d}x \int_0^{L/2} \mathrm{d}y \frac{\partial A^i}{\partial t} \cdot \nabla \times A.$$ \hspace{1cm} (6)

Then, since $\nabla \times A = B$ and $B = \mu_0(H + M)$, the signal

$$I_1 = -\frac{1}{V_0} \int_0^{2a} \mathrm{d}z \int_0^{L/2} \mathrm{d}x \int_0^{L/2} \mathrm{d}y \frac{\partial M^i}{\partial t}$$ \hspace{1cm} (7)

where here the time average is indicated by the bar. This final equation of the formalism is manifestly gauge invariant. The resonance of the ferromagnetic layer is contained in $\chi_r(t)$, the dynamic susceptibility, and

$$M_r(r, t) = \int \mathrm{d}r' \int \mathrm{d}t' \chi_r(r - r'; t - t') H_r(r', t'),$$ \hspace{1cm} (8)

where $i = x, y, z$. The simple expressions in equations (7) and (8) are a principal theoretical result presented here and have an obvious interpretation in terms of the magnetic energy. They demonstrate that Josephson junction magnetic spectroscopy measures very directly the magnetic susceptibility correlation function $\chi_r(r - r'; t - t')$, and all the other excitations to which that couples, in much the same manner as does neutron scattering. As will be seen below, the advantage is that this technique couples preferentially to small $q$-wave vectors.

Since $H_r(r, t)$ is periodic in time, the time convolution equation (8) reduces to a product and if the susceptibility is sufficiently local then

$$M_r(r, \omega_0) = (\chi_r'(\omega_0) + i\chi_r''(\omega_0)) H_r(r, \omega_0)$$ \hspace{1cm} (9)

in the usual complex notion. If rather the response is non-local then

$$M_r(k, \omega_0) = (\chi_r'(k, \omega_0) + i\chi_r''(k, \omega_0)) H_r(k, \omega_0)$$ \hspace{1cm} (10)

and it requires a Fourier expansion of the spatially dependent $H_r(r, \omega_0)$. Finally, when $\omega_0 \approx \xi_R$, as for the lowest voltage experimental signals, the linear response approximation is not strictly applicable and high harmonics of $\omega_0$ must be accounted for. Similar expressions apply but now, in particular, ‘half-harmonic’ signals occur since the super-current contains higher harmonics and can, corresponding to the second harmonic, excite a resonance at $\omega_0$ when $(\omega_0/2) = \omega_0$.

Now all that is required is a single integration of equation (4) but this is not a simple task. It is trivial to verify by differentiation that such an integral is $H_p = -(J_z/k) \cos(kx + \omega t)y$. However, this does not satisfy the boundary condition that the current density is zero outside the square junction region.

With the present gauge $A = A_z$ it is necessary to find a solution of Laplace’s equation $\nabla^2 A_z = 0$ such that $H = H_p + H_z$, where $H_p = \nabla \times A$ is such that $S$ satisfies equation (4) inside the square but has $\nabla \times H = 0$ outside. A little reflection suggests there are two contributions to $H_z$, which must be accounted for. First, in general, reflecting the even part $J_z \cos(kx + \omega t)$ of the current density there is a net oscillating super-current $I_z(H)$ that causes a circulation about the $z$-direction, and second, associated with the spatially odd part of $J_z \sin(kx + \omega t)$, there is a uniform component of the field in the $y$-direction. The first contribution is determined by considering the problem with $k = 0$, i.e. with a uniform current density $\mathcal{T}_z$. The solution is $H_z = \frac{1}{2} \mathcal{T}_z(x y - y k) \cos \omega t + \cdots$ where the ellipsis reflects the relatively small corrections for a square as compared with a circular cross section. In what follows this correction is ignored. The corresponding vector potential has $A_z = (1/4) \mathcal{T}_z (x^2 + y^2)$. For finite $H_z$ integrating the even part of the current density gives an average super-current density of $\mathcal{T}_z = J_z \sin(kL/2)/(kL/2)$. By inspection it is observed that $A_z = (1/4) \mathcal{T}_z (x^2 + y^2)$ satisfies $\nabla^2 A_z = 0$. The corresponding odd $H_z'$ correctly reduces to $H_z$ in the limit $k \to 0$. It is the case that $H_p$ and $H_p + H_z'$ implies a uniform oscillating field but one which diverges as $k \to 0$, whereas physically the even part of $H_p$ should be proportional to $k$ reflecting the Josephson screening of fields. That the tangential applied
field $H$ be continuous requires the current, induced by the even part of $H$, to be zero at the surface. The even part of $H_e$ is $-J_e\cos k_x \sin\omega t \hat{y}(\cos k_y/\phi_0)\sin\omega t \hat{y}$. The required even part of $H_i$ is now, $H_i = J_i \cos k_L \sin\omega t \hat{y}(\cos k_y/\phi_0)\sin\omega t \hat{y}$, which is equally divergent as $k \to 0$. The net result of integrating equation (4) is therefore

$$
H = \frac{J_e}{k} \left[ \cos(kx + \omega t) - \cos \frac{kL}{2} \sin\omega t \right] \hat{y} + \frac{1}{2} J_e \sin \frac{kL}{2} (-x \hat{y} - y \hat{k}) \cos\omega t + \cdots.
$$

(11)

Imagine that the magnetic layer is composed of a number of independent crystallites so that the response is local and equation (9) would apply. This local assumption also has the merit of being a useful illustration of the theory since it leads to a relatively simple prediction of the $H$ dependence of the signal which can be compared with experiment. This helps determine that the response is indeed local or extended. There are some complicated integrals involved in the evaluation of equation (7). The result is written, in closed form, as

$$
I_n = 2\pi \Phi_0 \frac{\Phi_i}{\Phi_0} \left[ F_x \chi''(\omega_0) + F_y \chi''(\omega_0) \right],
$$

(12)

where $\Phi_i = (2aL)B_{id} = (2aL)\mu_0 \Phi(0)/L$ is the flux due to the radio frequency field and where

$$
F_x = \frac{1}{48} \left[ \frac{I_x(B_0)}{I_c} \right]^2
$$

while

$$
F_y = \frac{2}{x^2} \left\{ 1 - \frac{x^2}{2} \sin^2 \frac{x}{2} \cos \frac{x}{2} - \frac{11}{24} + \frac{2}{x^2} \right\} \sin^2 \frac{x}{2},
$$

with $x = kL$, reflects the geometrical structure of the coupling. The equilibrium magnetization is along the $z$-axis, and the magnetic resonance signal is contained in $\chi''(\omega_0)$ and $\chi''(\omega_0)$, the Fourier transforms of the imaginary part of the susceptibility. The two functions $F_x$ and $F_y$ are plotted in figure 3. As required by symmetry $F_x = F_y$ when $k = 0$. The $\chi''(\omega_0)$ response near the first zero in $I_c(B_0) = I_c$ is about four times the maximum response to $\chi''(\omega_0)$. While there is a clear reflection of the Fraunhofer diffraction pattern in the $\chi''(\omega_0)$ response, one $1/k^2$, i.e. $1/H^2$, response dominates that of $\chi''(\omega_0)$ with only modulations due to diffraction effects. With a single flux quantum threading through the junction the $F_x$ is zero reflecting the net absence of a circulating current. In contrast $F_y$ is a maximum since the current is odd and the junction constitutes a small flat solenoid carrying the critical current density and hence a maximum field internal to the junction. In figure 3 we also report the amplitude of the FMR signal (markers) as a function of the applied magnetic field. The experimental data follow the $F_x$ coupling function.

Given a static magnetization $M = M_z \hat{k}$ the magnetic susceptibility might be approximated by,

$$
\chi''(k, \omega) = \chi''(k, \omega) \approx \gamma_s M_z \sum_{\pm} \frac{(\frac{4}{3})}{(\omega_0 + \alpha k^2 + (\frac{4}{3})^2 + \omega)^2}
$$

(13)

where $\tau$ is the relaxation time, $\omega_0$ the frequency of the FMR mode and $\alpha$ the spin-wave stiffness. Here $\gamma_s = \mu_B/2\mu_0$ with $\mu_0$ the Bohr magneton. We have performed simulations of the experimental signal using equations (12) and (13) with $\alpha = 0$, using $\omega_0$ and $M_z$ as parameters. Some of these are shown in figure 4. The signals here are believed to correspond to a half-harmonic signal and indeed the value of $M_z$ is smaller than that used to fit the regular FMR signal reported briefly elsewhere [2]. The frequency of that mode is given by Kittel’s formula $\omega_0 = \gamma_s \sqrt{(H_K - 4\pi M_S)^2 - H^2}$ where the anisotropy field $H_K$ and saturation magnetization $M_S$ are taken from separate measurements of these quantities and have a strength corresponding $M_z \approx M_S$, i.e. the observed signals are entirely consistent with the coupling of the FMR resonance to the superconductivity via the AB phase.

Kittel’s formula predicts that $\omega_0$ decreases with increasing $H$ whereas the half-harmonic signal of figure 4 actually increases. This leads us to believe that this signal corresponds to a finite value of $\alpha$ and $k$ in equation (13). That the spin-wave

Figure 3. In blue is $F_x$ and in pink $F_y$. These curves satisfy the evident requirement that the coefficient of the $x$ and $y$ responses be the same when $B = 0$.

Figure 4. Dynamical resistance of a ferromagnetic Josephson junction showing ferromagnetic resonances with the first and second harmonics. Inset: the magnetic field dependence of the resonance at the second harmonic.
The electromagnetic nature of the response can be confirmed by a study of Shapiro steps [1, 12]. In the absence of a magnetic resonance mode, an applied radio frequency field gives rise to such steps. In order to account for these we write for the bias voltage across the junction as

\[ V = V_0 + v \cos \Omega t \]

where \( v \) and \( \Omega \) are the amplitude and frequency of the applied microwave field while the constant \( V_0 \), as usual, corresponds to a Josephson frequency \( \omega_J = 2eV_0/h \). The Josephson current is now \( J_c \) times

\[ \sin \left( kx + \omega t + \frac{2ev}{\hbar \Omega} \sin \Omega t \right). \]

Expanding this sine with the Jacobi–Anger identity gives

\[ J_0 \left( \frac{2ev}{\hbar \Omega} \right) S_0 + \sum_{n=1}^{\infty} \left( J_{2n} \left( \frac{2ev}{h \Omega} \right) S_{2n} + J_{2n-1} \left( \frac{2ev}{h \Omega} \right) S_{2n-1} \right) \]

with \( S_{2n} = \sum_{\pm} \sin(kx + \omega_0 t + 2n \Omega t) \) and \( S_{2n-1} = \sum_{\pm} \sin(kx + \omega_0 t + (2n - 1) \Omega t) \), where \( S_0 = S_{2n=0} \). Each term is of the form:

\[ J_n^c \sin[kx + \omega_0 t + n \Omega t] \rightarrow J_n^c \sin[kx + \omega_0 t] \ (14) \]

where, in well known fashion [1], \( J_n^c = J_c \omega_n \) involves the appropriate Bessel function \( J_n(2ev/h \Omega) \). For the voltage at which \( \omega_0 = 0 \) there is a direct contribution \( J_n^c \sin(\phi_0 + kx) \) to the average current which is equivalent to the zero voltage critical current step but displaced to \( V_n = n\hbar \Omega/2e \) and reduced from \( I_c(H) \) to \( I_c(H) \omega_n \). This corresponds to the principal Shapiro steps, shown in the top panel of figure 5. For the magnetic response, in the linear response regime, the theory developed without an applied microwave field can be adopted. All that is needed is to replace \( J_c \) by \( J_n^c \) and \( \omega_0 \) with \( \omega_0 \) in the appropriate places as described above.

We have investigated the Shapiro steps for different microwave power. The left part of figure 5 shows the appearance of the Shapiro steps in the current-bias characteristics for increasing microwave power. The amplitude of each step follows the appropriate Bessel function \( J_n(2ev/h \Omega) \), as expected. In the bottom panel of figure 5 is presented the dependence of the magnetization induced sidebands on the microwave power. To make these bands more evident, we show the dynamical resistance as a function of the voltage. It is clear from the data that the amplitude of the sideband resonances follows the Bessel function of the Shapiro steps, as expected from the theory described above. Indeed the Shapiro steps can be seen as a ‘replica’ of the critical current at finite bias and hence the side-band amplitude follows that of the steps. The two side-bands correspond to the two poles of the dynamical susceptibility equation (13).

In conclusion, we have described experiments and developed theory to demonstrate that the relative AB phase of the superconductors which comprise a Josephson junction couples to the magnetic order parameter of a ferromagnet. We have thereby performed an FMR experiment with a sensitivity which greatly exceeds that of conventional cavity FMR. Since the coupling is via the magnetic field it is not necessary to have the current pass through the magnetic material. It might be imagined that the magnetic layer is the top layer of a ferromagnetic-superconducting-insulator-superconducting structure in which the adjacent S layer has a thickness of the order of, or less than, the London penetration length. The possibility of measuring the dispersion of spin-wave excitations has also been investigated. Our method [2, 11, 13] of coupling superconductivity to magnetism measures directly the dynamic susceptibility \( \chi''(q, \omega) \) with an enhanced sensitivity for small wavevectors, complementary to neutron scattering.

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