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Ferromagnet/superconductor heterostructures in graphene

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ABSTRACT

We investigate the unusual features of the proximity effect in graphene-based ferromagnet/superconductor (F/S) heterostructures. We show that, for a weakly doped ferromagnetic graphene with an exchange energy h exceeding its Fermi energy E_F, the Andreev reflection of massless Dirac fermions at the FS interface is accompanied with a Klein tunneling through an exchange field p–n barrier between two spin-split conduction and valence subbands. This spin Andreev–Klein process leads to an anomalous exchange-field-induced increase of the amplitude of Andreev reflection which has significant consequences for the subgap conductance and the shot noise power of an SF junction. For $h > E_F$, we find that the exchange field enhances the subgap conductance to reach a maximum value which is greater than the corresponding value for non-ferromagnetic structure at bias voltages $eV < \Delta/\sqrt{2}$. We further demonstrate that, due to the spin Andreev–Klein process, the Josephson coupling in a ferromagnetic graphene Josephson junction can be long-ranged, and that it survives at very large exchange fields $h \gg E_F$. The possibility for transitions between 0 and π Josephson coupling by varying the doping of the ferromagnetic contact is also demonstrated.

1. Introduction

It is a common belief that singlet superconductivity and ferromagnetic ordering do not coexist because of their incompatible natures. However, their interplay in artificially fabricated ferromagnet/superconductor (F/S) hybrid structures gives rise to exciting features which have been at the focus of interest in the past decade [1–3]. The experimental realization of these features became possible only recently due to the progress in the fabrication of high quality nanostructures [4]. The key process is the superconducting conversion of subgap electron–hole excitations at an interface between a normal (N) metal and a superconductor, known as Andreev reflection (AR) [5]. Andreev reflection makes possible the induction of the superconducting correlations into an N metal which is coupled to a superconductor. Due to this proximity effect FS systems provide a unique opportunity for studying the interplay of induced superconductivity with ferromagnetism.

More recently, proximity-induced superconductivity [6,7] as well as spin polarization [8–10] have been experimentally realized in the recently discovered graphene, two-dimensional (2D) carbon atoms arranged in hexagonal lattice [11–13]. These experimental achievements, providing the possibility of fabricating high quality graphene FS structures, have motivated theoretical studies of the interplay between superconductivity and ferromagnetism in this new carbon-based material [14–17]. Graphene is a zero-gap semiconductor with its conical valence and conduction bands touching each other at the corners of hexagonal first Brillouin zone, known as Dirac points. The carrier type, (electron-like (n) or hole-like (p)) and density can be tuned by means of an electrical gate or doping of the underlying substrate. The charge carriers in graphene behave like 2D massless Dirac fermions with a pseudo-relativistic chiral property [18–20,12,13]. Currently intriguing properties of graphene, which arise from such a Dirac-like spectrum, have been the subject of intense studies [21–23].

Already, the peculiarities of the proximity effect in graphene NS hybrid structures have been studied [25–28]. It was demonstrated that unlike highly doped graphene or a degenerate N metal, for undoped graphene with $\Delta \gg E_F$ the dominant process is AR of an electron from the conduction band into a hole in the valence band in which the reflection angle (versus the normal to the NS interface) is inverted with respect to the incidence angle, making the Andreev process a specular reflection [25,28].

In this paper, we demonstrate the unusual aspects of the proximity effect in hybrid FS structures which are realized in graphene. We show that the properties of these structures are drastically different in the two regimes of the Fermi energy smaller and larger than the exchange energy. Although for $E_F > h$ the amplitude of AR at an FS interface is suppressed by h as in the case of a metallic

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FS junction [24], for $E_F < h$ it is enhanced by the exchange splitting. We find that the subgap Andreev conductance of a graphene FS junction increases for $h > E_F$, and can even be larger than the corresponding value for a graphene NS junction at sufficiently low bias voltages. We explain this unusual effect in terms of a spin Andreev–Klein process in which the superconducting conversion of massless Dirac electron–hole excitations occurs. We also analyze the consequences of this quasi-relativistic Andreev process on the shot noise power of an FS junction [16]. For a graphene SFS junction we find that the spin Andreev–Klein process leads to a long-range Josephson coupling which is slowly damped only at very high exchange energies $h \gg E_F$ [14,17]. In addition we show that the 0–π Josephson coupling transitions can be seen by changing E_F from a value at $E_F < h$ to a value at $E_F > h$ for a certain length of ferromagnetic graphene junction, which is possible by varying the doping of the ferromagnetic contact using a gate voltage.

2. Transport in graphene FS junctions

We consider a wide graphene FS junction normal to the x-axis with a ferromagnetic region for $x < 0$ and a highly doped superconducting region for $x > 0$. To model the system we adopt the Dirac–Bogoliubov–de Gennes (DBdG) equation [25] which relates spin-s ($s = \pm$) electron and spin-$\bar{s}$ hole excitations,

$$\begin{pmatrix} H_s - E_F & \Delta(x) \\ \Delta(x)^* & E_F - H_{\bar{s}} \end{pmatrix} \begin{pmatrix} u_s \\ v_{\bar{s}} \end{pmatrix} = \varepsilon \begin{pmatrix} u_s \\ v_{\bar{s}} \end{pmatrix}, \quad (1)$$

$$H_s = -i\hbar v_F(\sigma_x \partial_x + \sigma_y \partial_y) - U(x) - sh(x), \quad (2)$$

$$U(x) = U_0\theta(x), \qquad h(x) = h\theta(-x), \qquad \Delta(x) = \Delta\theta(x), \quad (3)$$

where σ_x and σ_y are Pauli matrices in the pseudospin space of the sublattices and $\theta(x)$ is the step function. The potential U_0 in the superconducting region is very large in comparison with all other energy scales because this region should be highly doped. We consider $\varepsilon, h, E_F > 0$ and set $\hbar v_F = k_B = 1$ (k_B is the Boltzmann constant).

At a given energy ε and transverse momentum q we have these eigenstates inside F,

$$\psi^{\pm}_{es} = \frac{e^{\pm i k_s x}}{\sqrt{\cos\phi_s}} \begin{pmatrix} e^{\mp i\phi_s} \\ \pm 1 \\ 0 \\ 0 \end{pmatrix}, \qquad \psi^{\pm}_{h\bar{s}} = \frac{e^{\pm i k'_{\bar{s}} x}}{\sqrt{\cos\phi'_{\bar{s}}}} \begin{pmatrix} 0 \\ 0 \\ e^{\mp i\phi'_{\bar{s}}} \\ \mp 1 \end{pmatrix}, \quad (4)$$

with propagation angles $\phi_s = \arcsin(q/(\varepsilon + E_F + sh))$ and $\phi'_{\bar{s}} = \arcsin(q/(\varepsilon - E_F - \bar{s}h))$ with respect to the normal to the interface and longitudinal momenta $k_s = (\varepsilon + E_F + sh)\cos\phi_s$ and $k'_{\bar{s}} = (\varepsilon - E_F - \bar{s}h)\cos\phi'_{\bar{s}}$. Inside the superconductor we are interested in the states decaying as $x \to \infty$ for subgap energies ($\varepsilon < \Delta$),

$$\psi^{\pm}_S = e^{i k_{S\pm} x} \begin{pmatrix} e^{\pm i\beta} \\ \pm e^{\pm i\beta} \\ 1 \\ \pm 1 \end{pmatrix}, \quad k_{S\pm} = i\Delta\sin\beta \pm U_0, \quad (5)$$

where $\beta = \arccos(\varepsilon/\Delta)$.

For an incident spin-s electron from the left to the FS interface with a subgap energy $\varepsilon \le \Delta$ the scattering state in the two regions can be written as

$$\Psi_{e\sigma} = \begin{cases} \psi^{+}_{e\sigma} + r_\sigma \psi^{-}_{e\sigma} + r_{A\sigma}\psi^{-}_{h\bar{\sigma}}, & x < 0; \\ t^{+}\psi^{+}_S + t^{-}\psi^{-}_S, & x > 0, \end{cases} \quad (6)$$

where r_s and r_{As} denote the amplitude of normal and Andreev reflections, respectively. Since for energies below the superconducting gap transmission into the superconductor is forbidden, we will have $|r_s|^2 = 1 - |r_{As}|^2$. Therefore, for transport properties (conductivity, noise power, etc.) it is sufficient to have only AR probability $R_{As} \equiv |r_{As}|^2$, which can be found from continuity condition at the interface as

$$R_{As} = \frac{2\cos\phi_s\cos\phi'_{\bar{s}}}{1 + \cos\phi_s\cos\phi'_{\bar{s}} + \cos 2\beta\sin\phi_s\sin\phi'_{\bar{s}}} \quad (7)$$

in which the reflected hole angle is given by

$$\frac{\sin\phi'_{\bar{\sigma}}}{\sin\phi_\sigma} = P_s = \frac{\varepsilon + E_F + sh}{\varepsilon - E_F + sh}. \quad (8)$$

Although in usual metal–superconductor interfaces we always have Andreev retro-reflection ($\phi' = -\phi$), Eq. (8) shows that for a graphene FS junction we have a variety of processes depending on the values of ε and h with respect to E_F as well as the spin of incident electrons. This essential difference, first pointed out by Beenakker for the special case of a graphene NS junction [25, 28], arises from the gapless two-band (conduction and valence) structure of graphene which gives the possibility of both intra-band and inter-band AR.

To further illustrate the effect we examine different situations using Eq. (8). Considering a spin-s electron, while $0 < \varepsilon + sh < E_F$, the reflected spin-$\bar{s}$ hole is from the same band (conduction or valence) and we have retro-type reflection $\mathrm{sgn}\,\phi' = -\mathrm{sgn}\,\phi$. By increasing $\varepsilon + sh$ the reflected hole angle is also increased, and when we reach the limit $0 < \varepsilon + sh = E_F$ it jumps from $-\pi/2$ to $\pi/2$. So for $\varepsilon + sh > E_F$ the reflection is specular type, and a further increase in $\varepsilon + sh$ decreases $\phi'_{\bar{s}}$ to reach ϕ_s for very large $\varepsilon + sh$. In all the above situations we have $0 < \varepsilon + sh$ and the AR is divergent ($|\phi'_{\bar{s}}| > |\phi_s|$). Due to this divergent behavior, for electrons with incidence angles above a critical value $\phi^c_s = \arcsin(1/|P_s|)$ the reflected hole is evanescent and thus such electrons do not contribute in the transport through the junction. On the other hand, for states with $\varepsilon + sh < 0$ (this situation can occur for spin-down electrons when $h > \varepsilon$), the reflection becomes convergent. In this case by decreasing $\varepsilon + sh$ the reflection angle increases continuously from $-\phi_s$ to ϕ_s, passing through 0 for $\varepsilon + sh = -E_F$. The different possibilities for AR are shown schematically in Fig. 1 separately for two spin directions.

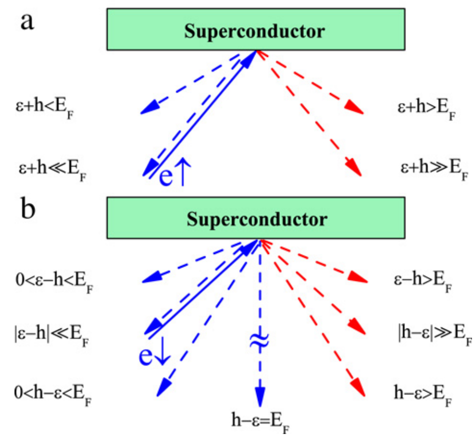

Fig. 1. Possible trajectories of an incident (a) up and (b) down electron (solid lines) and the Andreev reflected hole (dashed lines), as a function of the excitation energy ε and exchange field h compared to the Fermi energy E_F. The Andreev reflection can be retro (blue lines) or specular (red lines) for both spin directions depending on whether the reflected hole is from the same band as the incident electron or not. For a spin-s electron with incidence angle ϕ_s, when $\varepsilon + sh$ increases from 0 to $+\infty$ the reflection angle rotates clockwise from $-\phi_s$ to ϕ_s, jumping from $-\pi/2$ to $\pi/2$ at $\varepsilon + sh = E_F$. For spin-down electrons it is also possible that $\varepsilon - h$ varies from 0 to $-\infty$, which leads to the counterclockwise rotation of the reflection angle between $-\pi/2$ and $\pi/2$. (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this article.)