Scanning tunneling spectroscopy and Recent progress of high-temperature superconductors

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I. INTRODUCTION

In relativistic quantum mechanics spin–1/2 fermions are described by the solutions of the Dirac equation [1],

$$(i\gamma^\mu \partial_\mu - m)\psi = 0. \tag{1}$$

Here, γ^μ ($\mu = 0, 1, 2, 3$) is a set of 4×4 matrices satisfying the anti–commutation relations $\{\gamma^\mu, \gamma^\nu\} = 2g^{\mu\nu}$, where $g^{\mu\nu}$ represents the Minkowski metric and $\gamma_0\gamma_\mu\gamma_0 = \gamma_\mu^\dagger$. In general, the matrices γ^μ have complex elements, making Eq. (1) a set of coupled differential equations with complex coefficients.Thus, the general solution $\psi(x)$ of Eq. (1) representing the fermion field is a complex bi–spinor that is not an eigenstate of the charge conjugation (CC) operator. Since charge conjugation, $\psi \to \psi^*$, maps a particle into its anti–particle [2], a complex solution of Eq. (1) represents a fermion that has a distinct anti–fermion, with the same mass and spin but opposite charge and magnetic moment, as its counterpart. Therefore, the field of a relativistic fermion that coifncides with its own antiparticle, should it exist, is necessarily an eigenstate of CC that must be described by a real solution $\chi(x)$ of Eq. (1). Real solutions of the Dirac equation are possible, provided one can find a suitable representation of the matrices γ^μcharacterized by purely imaginary non–zero matrix elements. Such a representation of the γ^μ matrices that renders Eq. (1) purely real was found by E. Majorana in 1937 [3].

II. ANOTHER SECTION

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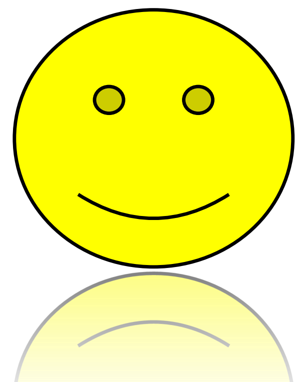

FIG. 1. This is a smiley face

A. A subsection

In 2D the mean field Hamiltonian for a spinless $(p_x + ip_y)$ superconductor (superfluid) is given by,

$$H_{2D}^p = \sum_p \xi_p c_p^\dagger c_p + \Delta_0 \sum_p \left[(p_x + ip_y) c_p^\dagger c_{-p}^\dagger + h.c. \right], \tag{2}$$

where $\xi_p = p^2/2m - \epsilon_F$, with ϵ_F the Fermi energy, and spin indices are omitted because the system is considered spinless (or spin-polarized).

B. Another subsection

Next, we assume that only a few bands are occupied, and that the low–energy subspace is defined by the eigenstates satisfying the condition $\epsilon_{\mathbf{n}} < \epsilon_{\max}$, where the cutoff energy $\epsilon_{\max}$ is typically of the order 100meV. Using this low-energy basis, the matrix elements of the total Hamiltonian can be written

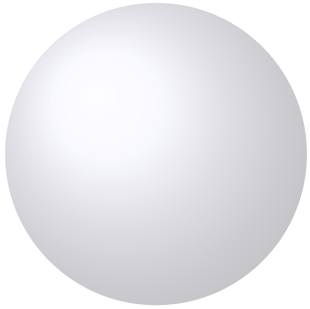

FIG. 2. This is a sphere.

explicitly. The matrix elements of the SOI Hamiltonian are

$$\langle\psi_{\boldsymbol{n}\sigma}|H_{\rm SOI}|\psi_{\boldsymbol{n}'\sigma'}\rangle = \alpha\delta_{n_z n_z'}\left\{q_{n_x n_x'}(i\sigma_y)_{\sigma\sigma'}\delta_{n_y n_y'} - q_{n_y n_y'}(i\sigma_x)_{\sigma\sigma'}\delta_{n_x n_x'}\right\}, \quad (3)$$

where

$$q_{n_\lambda n_\lambda'} = \frac{1-(-1)^{n_\lambda+n_\lambda'}}{N_\lambda+1}\frac{\sin\frac{\pi n_\lambda}{N_\lambda+1}\sin\frac{\pi n_\lambda'}{N_\lambda+1}}{\cos\frac{\pi n_\lambda}{N_\lambda+1}-\cos\frac{\pi n_\lambda'}{N_\lambda+1}}. \quad (4)$$

The first term in Eq. (3) represents the intra–band Rashba spin–orbit interaction, while the second term couples different confinement–induced bands. Similarly, assuming that the Zeeman splitting Γ is generated by a magnetic field oriented along the wire (i.e., along the x-axis), $\Gamma = g^*\mu_B B_x/2$, where g^* is the effective g–factor for the SM nanowire, the matrix elements for the corresponding term in the Hamiltonian are

$$\langle\psi_{\boldsymbol{n}\sigma}|H_{\rm Z}|\psi_{\boldsymbol{n}'\sigma'}\rangle = \Gamma\delta_{\boldsymbol{n}\boldsymbol{n}'}\delta_{\bar{\sigma}\sigma'}, \quad (5)$$

where $\bar{\sigma} = -\sigma$. Adding together these contributions, the effective Hamiltonian describing the low–energy physics of the semiconductor nanowire in the presence of a Zeeman field becomes $H_{\boldsymbol{n}\boldsymbol{n}'} = \langle\psi_{\boldsymbol{n}}|H_{\rm SM}+H_{\rm Z}|\psi_{\boldsymbol{n}}'\rangle$, with $\psi_{\boldsymbol{n}}$ representing the spinor $(\psi_{\boldsymbol{n}\uparrow}, \psi_{\boldsymbol{n}\downarrow})$. Explicitly, we have

$$H_{\boldsymbol{n}\boldsymbol{n}'} = \left[\epsilon_{\boldsymbol{n}} + \Gamma\sigma_x\right]\delta_{\boldsymbol{n}\boldsymbol{n}'} + i\alpha\delta_{n_z n_z'}\left[q_{n_x n_x'}\sigma_y\delta_{n_y n_y'} - q_{n_y n_y'}\sigma_x\delta_{n_x n_x'}\right], \quad (6)$$

where $\boldsymbol{n} = (n_x, n_y, n_z)$ and $q_{n_\lambda n_\lambda'}$ by Eq. (4). Note that a similar effective low–energy Hamiltonian can be written for an infinite quasi–1D wire. In the limit $L_x \to \infty$, the wave vector k_x becomes a good quantum number and we have

$$H_{\boldsymbol{n}\boldsymbol{n}'}(k_x) = [\epsilon_{\boldsymbol{n}}(k_x)+\alpha_R k_x\sigma_y+\Gamma\sigma_x]\delta_{\boldsymbol{n}\boldsymbol{n}'} - i\alpha q_{n_y n_y'}\sigma_x\delta_{n_z n_z'}, \quad (7)$$

where $\alpha_R = \alpha a$ and $\boldsymbol{n} = (n_y, n_z)$ labels the confinement–induced bands with energy

$$\epsilon_{\boldsymbol{n}}(k_x) = \frac{\hbar^2 k_x^2}{2m^*} - 2t_0\left(\cos\frac{\pi n_y}{N_y+1} + \cos\frac{\pi n_z}{N_z+1} - 2\right) - \mu. \quad (8)$$

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[2] M. Gell-Mann and A. Pais, Phys. Rev. **97**, 1387 (1955).
[3] E. Majorana, Nuovo Cimento **5**, 171 (1937).