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LETTER TO THE EDITOR

Spin analogue of the controlled Josephson charge current

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Abstract
We propose a controlled Josephson spin current across the junction of two non-centrosymmetric superconductors like $CePt_3Si$. The Josephson spin current arises due to the direction-dependent tunnelling matrix element and different momentum-dependent phases of the triplet components of the gap function. Its modulation with the angle ξ between the non-centrosymmetric axes of two superconductors is proportional to $\sin \xi$. This particular dependence on ξ may find application of the proposed set-up in making a *Josephson spin switch*.

Traditionally, Josephson junctions [1, 2] in superconductors draw interest both scientifically and in their applicability in making devices. With no exception, they have also been studied in unconventional superconductors like spin-singlet cuprate [3] and spin-triplet Sr_2RuO_4 [4] superconductors. However, no Josephson junction between nonmagnetic superconductors is known to generate spin-polarized current. The purpose of this letter is to theoretically show that the direction-dependent tunnelling matrix element across the junction between two recently discovered [5] non-centrosymmetric superconductors like $CePt_3Si$ leads to tunnelling of both spin-singlet and spin-triplet Cooper pairs. As a consequence, a non-vanishing spin-Josephson current is viable along with the usual charge-Josephson current. This *novel* spin-Josephson current depends on the relative angle ξ between the axes of non-centrosymmetry $\hat{n}_L$ and $\hat{n}_R$ in the left and right side of the junction respectively. This angular dependence may be used to make a *Josephson spin switch*.

The normal state Hamiltonian [6, 7] for the electrons in a band of a lattice without inversion symmetry is

$$H_0 = \sum_{\mathbf{k},s} \xi_{\mathbf{k}} c_{\mathbf{k}s}^\dagger c_{\mathbf{k}s} + \sum_{\mathbf{k},s,s'} \mathbf{g}_{\mathbf{k}} \cdot \boldsymbol{\sigma}_{ss'} c_{\mathbf{k}s}^\dagger c_{\mathbf{k}s'}, \tag{1}$$

where electrons with momentum $\mathbf{k}$ and spin s ($=\uparrow$ or $\downarrow$) are created (annihilated) by the operators $c_{\mathbf{k}s}^\dagger$ ($c_{\mathbf{k}s}$), and $\xi_{\mathbf{k}}$ is the band energy measured from the Fermi energy ϵ_F. The second term in the Hamiltonian (1) breaks parity symmetry as $\mathbf{g}_{-\mathbf{k}} = -\mathbf{g}_{\mathbf{k}}$ for a non-centrosymmetric system. For a system like the heavy fermion compound $CePt_3Si$ which has

a layered structure, H_0 is considered to be two-dimensional. For such a system of electrons with band mass m, $\xi_{\mathbf{k}} = \frac{k^2}{2m} - \epsilon_F$ and $g_{\mathbf{k}} = \alpha\boldsymbol{\eta}_{\mathbf{k}}$, where $\boldsymbol{\eta}_k = \hat{n} \times \mathbf{k}$, i.e., the spin–orbit interaction is of Rashba type [8] and α is the Rashba parameter. Here $\hat{n}$ represents the axis of non-centrosymmetry which is perpendicular to the plane of the system. Due to the breaking down of the parity, spin degeneracy of the band is lifted; by diagonalizing H_0, one finds two spin-split bands with energies $\xi_{\mathbf{k}\lambda} = \xi_k + \lambda\alpha|\mathbf{k}|$, where $\lambda = \pm$ describes the helicity of the spin-split bands. Therefore in the diagonalized basis H_0 (1) becomes $H_0 = \sum_{\mathbf{k},\lambda=\pm} \xi_{\mathbf{k}\lambda}\tilde{c}^\dagger_{\mathbf{k}\lambda}\tilde{c}_{\mathbf{k}\lambda}$, where $\tilde{c}_{\mathbf{k}\lambda} = (c_{\mathbf{k}\uparrow} - i\lambda\exp(i\phi_{\mathbf{k}})c_{\mathbf{k}\downarrow})/\sqrt{2}$ is the electron destruction operator and $\tilde{c}^\dagger_{\mathbf{k}\lambda} = (c^\dagger_{\mathbf{k}\uparrow} + i\lambda\exp(-i\phi_{\mathbf{k}})c^\dagger_{\mathbf{k}\downarrow})/\sqrt{2}$ is the electron creation operator in band λ with momentum $\mathbf{k}$ whose orientation with $\hat{x}$-axis is $\phi_{\mathbf{k}}$. The density of electronic states at the Fermi energy in these bands may be found as $\nu_\lambda = \frac{m}{2\pi}(1 - \lambda m\alpha/\sqrt{k_F^2 + m^2\alpha^2})$, where $k_F = \sqrt{2m\epsilon_F}$ is the Fermi momentum.

Band structure calculation [9] on CePt$_3$Si reveals that the energy difference between the two spin-split bands near k_F is 50–200 meV, which is much larger than the superconducting critical temperature, $k_B T_c \approx 0.06$ meV [5]. The formation of Cooper pairing between electrons in different spin-split bands may thus be ignored. The Hamiltonian for these superconductors may then be written as

$$H_1 = \sum_{\mathbf{k},\lambda=\pm} \left[\xi_{\mathbf{k}\lambda}\tilde{c}^\dagger_{\mathbf{k}\lambda}\tilde{c}_{\mathbf{k}\lambda} + \left(\Delta_{\mathbf{k}\lambda}\tilde{c}^\dagger_{\mathbf{k}\lambda}\tilde{c}^\dagger_{-\mathbf{k}\lambda} + \text{h.c.} \right) \right], \tag{2}$$

where Cooper pairs are only between intraband electrons. Therefore the normal and anomalous Green's functions are obtained respectively as $\mathcal{G}_\lambda(\mathbf{k}, i\epsilon_n) = -(i\epsilon_n + \xi_{\mathbf{k}\lambda})/(\epsilon_n^2 + \xi_{\mathbf{k}\lambda}^2 + |\Delta_{\mathbf{k}\lambda}|^2)$ and $\mathcal{F}_\lambda(\mathbf{k}, i\epsilon_n) = \Delta_{\mathbf{k}\lambda}/(\epsilon_n^2 + \xi_{\mathbf{k}\lambda}^2 + |\Delta_{\mathbf{k}\lambda}|^2)$, where ϵ_n is the fermionic Matsubara frequency [13]. The superconducting order parameter $\Delta_{\mathbf{k}\lambda}$ obeys the symmetry [9, 10]: $\Delta_{-\mathbf{k}\lambda} = -\Delta_{\mathbf{k}\lambda}$. We consider $\Delta_{\mathbf{k}\lambda} = \tilde{\Delta}_{k\lambda}\Lambda_{\mathbf{k}}$, i.e., the angular dependences of $\mathbf{k}$ on the order parameters of the two bands are assumed to be same. Apart from the overall phase rigidity angle Θ of the superconductor, there may be a relative phase difference θ between the two bands: $\tilde{\Delta}_{k+} = |\tilde{\Delta}_{k+}|e^{i\Theta}$ and $\tilde{\Delta}_{k-} = |\tilde{\Delta}_{k-}|e^{i(\Theta+\theta)}$. Reverting H_1 (2) to a spin-up (-down) basis, we find

$$H_1 = \sum_{\mathbf{k}} \Psi^\dagger_{\mathbf{k}} \begin{pmatrix} \xi_{\mathbf{k}} & \Delta_{\mathbf{k},\uparrow\uparrow} & \Gamma_{\mathbf{k},R} & \Delta_{\mathbf{k},\uparrow\downarrow} \\ \Delta^*_{\mathbf{k},\uparrow\uparrow} & -\xi_{\mathbf{k}} & \Delta^*_{\mathbf{k},\downarrow\uparrow} & \Gamma^*_{\mathbf{k},R} \\ \Gamma^*_{\mathbf{k},R} & \Delta_{\mathbf{k},\downarrow\uparrow} & \xi_{\mathbf{k}} & \Delta_{\mathbf{k},\downarrow\downarrow} \\ \Delta^*_{\mathbf{k},\uparrow\downarrow} & \Gamma_{\mathbf{k},R} & \Delta^*_{\mathbf{k},\downarrow\downarrow} & -\xi_{\mathbf{k}} \end{pmatrix} \Psi_{\mathbf{k}}, \tag{3}$$

where $\Delta_{\mathbf{k},\uparrow\uparrow} = \frac{1}{2}(\tilde{\Delta}_{k+} + \tilde{\Delta}_{k-})\Lambda_{\mathbf{k}}$, $\Delta_{\mathbf{k},\downarrow\downarrow} = \frac{1}{2}\exp[2i\phi_{\mathbf{k}}](\tilde{\Delta}_{k+} + \tilde{\Delta}_{k-})\Lambda_{\mathbf{k}}$, and $\Delta_{\mathbf{k},\uparrow\downarrow} = -\Delta_{\mathbf{k},\downarrow\uparrow} = \frac{i}{2}\exp[i\phi_{\mathbf{k}}](\tilde{\Delta}_{k+} - \tilde{\Delta}_{k-})\Lambda_{\mathbf{k}}$ are the different components of pairing potential $\Delta_{\mathbf{k},ss'}$ between electrons with spins s and s', and $\Gamma_{\mathbf{k},R} = i\alpha|\mathbf{k}|\exp[-i\phi_{\mathbf{k}}]$ is the Rashba spin–orbit coupling potential. The Hamiltonian (3) has been expressed in the basis such that $\Psi^\dagger_{\mathbf{k}} = (c^\dagger_{\mathbf{k}\uparrow}, c_{-\mathbf{k}\uparrow}, c^\dagger_{\mathbf{k}\downarrow}, c_{-\mathbf{k}\downarrow})$. Since $\Delta_{\mathbf{k},\uparrow\downarrow} = -\Delta_{\mathbf{k},\downarrow\uparrow}$, there is no triplet component with zero projection along the spin-quantization direction; pairing between electrons with unequal spins entirely gives rise to the singlet component. If we choose $\Lambda_{\mathbf{k}} = -i\exp[-i\phi_{\mathbf{k}}]$, the triplet component of the pairing may be expressed as $\hat{\Delta}_{\mathbf{k},T} = (d_{\mathbf{k}} \cdot \boldsymbol{\sigma})i\sigma_y$ with $d_{\mathbf{k}} = \frac{1}{2|\mathbf{k}|}(\tilde{\Delta}_{k+} + \tilde{\Delta}_{k-})\boldsymbol{\eta}_{\mathbf{k}}$. In other words, the only stable [7] spin triplet component whose $d_{\mathbf{k}} \parallel \boldsymbol{\eta}_{\mathbf{k}}$ corresponds to this $\Lambda_{\mathbf{k}}$. Therefore $\Delta_{\mathbf{k},\uparrow\uparrow}$ and $\Delta_{\mathbf{k},\downarrow\downarrow}$ will have equal but opposite momentum-dependent phase. Accordingly the singlet component of the pairing potential becomes $\Delta_{\mathbf{k},S} = \frac{1}{2}(\tilde{\Delta}_{k+} - \tilde{\Delta}_{k-})$. If $|\tilde{\Delta}_{k+}| \neq |\tilde{\Delta}_{k-}|$, the admixture [6, 7] of singlet and triplet pairing takes place. Recent observation [11] of Josephson current in the junction of CePt$_3$Si and an s-wave superconductor suggests the existence of a spin-singlet order parameter, while larger [5] upper critical field H_{c2} seems to suggest that spin-triplet pairing occurs. One normally assumes the
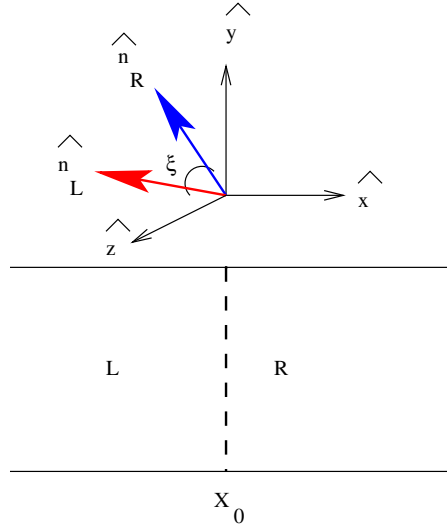
Figure 1. Schematic picture of two non-centrosymmetric superconductors, denoted by L and R, coupled at $x = x_0$. Their orientations are characterized by the axes of non-centrosymmetry, $\hat{n}_L$ and $\hat{n}_R$.

(This figure is in colour only in the electronic version)

superconducting order parameter to be independent of the magnitude of momentum, i.e., $\tilde{\Delta}_{k\lambda}$ is k-independent. The anomalous Green's function $\mathcal{F}_{ss'}(\mathbf{k}, \mathrm{i}\epsilon_n)$ in the up–down basis is related to $\mathcal{F}_\lambda$ as $\mathcal{F}_{\uparrow\uparrow} = \frac{1}{2}(\mathcal{F}_+ + \mathcal{F}_-)$, $\mathcal{F}_{\downarrow\downarrow} = \frac{1}{2}e^{2\mathrm{i}\phi_\mathbf{k}}(\mathcal{F}_+ + \mathcal{F}_-)$, and $\mathcal{F}_{\uparrow\downarrow} = -\mathcal{F}_{\downarrow\uparrow} = \frac{\mathrm{i}}{2}e^{\mathrm{i}\phi_\mathbf{k}}(\mathcal{F}_+ - \mathcal{F}_-)$. Apart from the applicable momentum-dependent phases, the triplet (singlet) component of $\mathcal{F}_{ss'}$ is the addition (subtraction) of the anomalous Green's function of the two spin-split bands.

We now consider Josephson tunnelling between two such non-centrosymmetric superconductors as depicted in figure 1. The Hamiltonian for the system then reads $H = H_L + H_R + H_T$, where H_L and H_R are the bulk Hamiltonian of the left and right side of the junction respectively and H_T describes the tunnelling between these two sides. The Hamiltonian H_L is described by H_1 (2) and H_R is also defined equivalently with the change in the notation of momentum $\mathbf{k} \rightarrow \mathbf{p}$ to distinguish each side of the junction. Further, each parameter in the left (right) is denoted by superscript or subscript L (R). The tunnelling Hamiltonian reads $H_T = \sum_{\mathbf{kp},s}(T_{\mathbf{kp}}c_{\mathbf{k}s}^\dagger c_{\mathbf{p}s} + T_{\mathbf{kp}}^* c_{\mathbf{p}s}^\dagger c_{\mathbf{k}s})$, where $T_{\mathbf{kp}}$ is the tunnelling matrix element for an electron with momentum $\mathbf{p}$ to tunnel from the right side to the left side with momentum $\mathbf{k}$. Time-reversal symmetry of H_T suggests $T_{-\mathbf{k},-\mathbf{p}} = T_{\mathbf{kp}}^*$.

In a charge-tunnelling process, the number of electrons in either side of the junction changes with time, e.g., the rate of change in the number of electrons of the right side of the junction $N_R = \sum_{\mathbf{p},s} c_{\mathbf{p}s}^\dagger c_{\mathbf{p}s}$ reads as $\dot{N}_R = \mathrm{i}[H_T, N_R]$. Therefore the tunnelling charge current becomes $I_c(t) = -e\langle \dot{N}_R(t)\rangle$: one part corresponds to quasi-particle tunnelling and the other part is due to the process of Josephson tunnelling [12, 13], i.e., tunnelling of the Cooper pairs. We are herewith interested in the Josephson charge current defined as $I_J^c(t) = 2e\,\mathrm{Im}[\exp(-2\mathrm{i}eVt)\Phi_{\mathrm{ret}}^c(eV)]$ at time t and for an applied bias voltage V across the junction, where $\Phi^c(\mathrm{i}\omega_m) = -\int_0^\beta \mathrm{d}\tau\, e^{\mathrm{i}\omega_m\tau}\langle T_\tau A(\tau)A(0)\rangle$ with imaginary time τ, $A(\tau) = \sum_{\mathbf{kp}s} T_{\mathbf{kp}}c_{\mathbf{k}s}^\dagger(\tau)c_{\mathbf{p}s}(\tau')$, the bosonic Matsubara frequency [13] ω_m, and inverse temperature β. Instead of the axes of non-centrosymmetry of the two sides of the junction being parallel,

general axes $\hat{n}_L$ and $\hat{n}_R$ will lead to $\hat{\boldsymbol{\eta}}_k = \hat{n}_L \times \hat{\mathbf{k}}$ and $\hat{\boldsymbol{\eta}}_p = \hat{n}_R \times \hat{\mathbf{p}}$. In that case, Josephson tunnelling occurs in both the singlet and triplet channels. Separating these channels, we find

$$\Phi^c(i\omega_m) = -\sum_{i\epsilon_n, \mathbf{kp}} \frac{|T_{\mathbf{kp}}|^2}{2\beta} \left[\chi_S + (\hat{\boldsymbol{\eta}}_k \cdot \hat{\boldsymbol{\eta}}_p) \chi_T \right], \tag{4}$$

where

$$\chi_S = \sum_{\lambda,\lambda'} \lambda\lambda' \Lambda_{\mathbf{k}} \Lambda_{\mathbf{p}}^* \mathcal{F}_\lambda^\dagger(\mathbf{k}, i\epsilon_n) \mathcal{F}_{\lambda'}(\mathbf{p}, i\epsilon_n - i\omega_m) \tag{5}$$

and

$$\chi_T = \sum_{\lambda,\lambda'} \Lambda_{\mathbf{k}} \Lambda_{\mathbf{p}}^* \mathcal{F}_\lambda^\dagger(\mathbf{k}, i\epsilon_n) \mathcal{F}_{\lambda'}(\mathbf{p}, i\epsilon_n - i\omega_m) \tag{6}$$

represent singlet and triplet contributions respectively. While the direction-independent electronic tunnelling probability $|T_{\mathbf{kp}}|^2 = |T|^2$ leads to tunnelling of the singlet component only, the triplet components of Cooper pairs tunnel due to the realistic direction-dependent [14] $|T_{\mathbf{kp}}|^2$. This triplet part in the charge Josephson tunnelling process may effectively provide spin-Josephson tunnelling. The rate of change of spin $S_z = \sum_{\mathbf{p}} [c_{\mathbf{p}\uparrow}^\dagger c_{\mathbf{p}\uparrow} - c_{\mathbf{p}\downarrow}^\dagger c_{\mathbf{p}\downarrow}]$ due to tunnelling, i.e., $\dot{S}_z = i[H_T, S_z]$, gives rise to the spin-Josephson tunnelling current $I_J^s(t) = -2\,\mathrm{Im}[\exp(-2ieVt)\Phi_{\mathrm{ret}}^s(eV)]$, where $\Phi^s(i\omega_m) = -\int_0^\beta d\tau\, e^{i\omega_m\tau}\langle T_\tau B(\tau)A(0)\rangle$ with $B(\tau) = \sum_{\mathbf{kp}} T_{\mathbf{kp}} [c_{\mathbf{k}\uparrow}^\dagger(\tau)c_{\mathbf{p}\uparrow}(\tau) - c_{\mathbf{k}\downarrow}^\dagger(\tau)c_{\mathbf{p}\downarrow}(\tau)]$. We thus find

$$\Phi^s(i\omega_m) = i\sum_{i\epsilon_n, \mathbf{kp}} \frac{|T_{\mathbf{kp}}|^2}{2\beta} \sum_j (\hat{\boldsymbol{\eta}}_k \times \hat{\boldsymbol{\eta}}_p)_j \chi_T, \tag{7}$$

where $(\hat{\boldsymbol{\eta}}_k \times \hat{\boldsymbol{\eta}}_p)_j$ represents the jth spatial component of $\hat{\boldsymbol{\eta}}_k \times \hat{\boldsymbol{\eta}}_p$.

We assume that the non-centrosymmetric axes are parallel to the interface of the junction as shown in figure 1. The probability of tunnelling [14] will be greatest along the direction transverse to the interface. In terms of normal state electronic tunnelling conductance $G_N = I_N/V$ with I_N being normal state tunnelling current, we find the dc ($V = 0$) charge and spin Josephson currents:

$$I_J^c = \left(\frac{G_N}{e\mathcal{F}}\right) \left[\sin\psi(g_+\mathcal{A}_1 + g_-\mathcal{A}_2) + \cos\psi(g_+\mathcal{A}_3 + g_-\mathcal{A}_4)\right] \tag{8}$$

$$I_J^s = \left(\frac{G_N}{e^2\mathcal{F}}\right) \delta\kappa \sin\xi \left[\cos\psi\,(\mathcal{A}_1 + \mathcal{A}_2) - \sin\psi\,(\mathcal{A}_3 + \mathcal{A}_4)\right], \tag{9}$$

where $\psi = \Theta_L - \Theta_R$, $g_\pm = \delta\cos\xi \pm 1$, $\delta \in (0, 1)$ is a parameter depending on the model[1] of the tunnelling matrix element, i.e., the angular dependence of $T_{\mathbf{kp}}$, $\cos\xi = \hat{n}_L \cdot \hat{n}_R$, where both $\hat{n}_L$ and $\hat{n}_R$ are in the plane of the interface, and κ is the projection of $\hat{n}_L \times \hat{n}_R$ along the direction perpendicular to the interface. Further, $\mathcal{F} = \pi \sum_{\lambda\lambda'} \nu_\lambda^L \nu_{\lambda'}^R (1 + \delta\lambda\lambda' \cos\xi)$ describes the dependence of normal state tunnelling current on δ, $\mathcal{A}_1 = \Gamma_{++} + \Gamma_{--}\cos(\theta_L - \theta_R)$, $\mathcal{A}_2 = \Gamma_{-+}\cos\theta_L + \Gamma_{+-}\cos\theta_R$, $\mathcal{A}_3 = \Gamma_{--}\sin(\theta_L - \theta_R)$, and $\mathcal{A}_4 = \Gamma_{-+}\sin\theta_L - \Gamma_{+-}\sin\theta_R$ with

$$\Gamma_{\lambda\lambda'} = \pi \nu_\lambda^L \nu_{\lambda'}^R \frac{|\tilde{\Delta}_\lambda^L||\tilde{\Delta}_{\lambda'}^R|}{|\tilde{\Delta}_\lambda^L| + |\tilde{\Delta}_{\lambda'}^R|} K\left(\frac{|\tilde{\Delta}_\lambda^L| - |\tilde{\Delta}_{\lambda'}^R|}{|\tilde{\Delta}_\lambda^L| + |\tilde{\Delta}_{\lambda'}^R|}\right) \tag{10}$$

at zero temperature, where K is an elliptic function of the first kind.

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[1] For example, $\delta \approx 0.4$ for $|T_{\mathbf{kp}}|^2 \sim |T|^2 \hat{k}_\perp \hat{p}_\perp \Theta(\hat{k}_\perp \hat{p}_\perp)$ as suggested in [15], where Θ represents the Heaviside function and subscript $\perp$ represents the direction of momentum along the direction transverse to the junction.

The charge (8) and spin (9) dc Josephson currents depend on both the sine and cosine of the global phase difference ψ between the two superconductors. Instead of a continuous relative phase difference between the superconductors in each spin-split band, there is a possibility of phase locking such that $\theta_L, \theta_R = n\pi$ ($n \in Z$). In this case, $\mathcal{A}_3 = \mathcal{A}_4 = 0$ and hence $I_J^c(V = 0) \propto \sin\psi$ and $I_J^s(V = 0) \propto \cos\psi$. This situation is also true for $\theta_L = \theta_R$, even in the absence of their locking at $n\pi$. The critical charge- and spin-Josephson currents are modulated with the angle ξ.

For general values of θ_L and θ_R, the charge (8) and spin (9) Josephson currents may be parameterized as $I_J^c = J_1^c \sin(\psi + \chi_1) + J_2^c \sin(\psi + \chi_2)$ and $I_J^s = J^s \cos(\psi + \chi_1)$, where $J_1^c(e\mathcal{F}/G_N) = \mathcal{C}\delta\cos\xi$, $J_2^c(e\mathcal{F}/G_N) = \mathcal{D}$, $J^s(e^2\mathcal{F}/G_N) = \mathcal{C}\delta\kappa\sin\xi$, $\chi_1 = \tan^{-1}(\frac{\mathcal{A}_3+\mathcal{A}_4}{\mathcal{A}_1+\mathcal{A}_2})$ and $\chi_2 = \tan^{-1}(\frac{\mathcal{A}_3-\mathcal{A}_4}{\mathcal{A}_1-\mathcal{A}_2})$ with $\mathcal{C} = [(\mathcal{A}_1+\mathcal{A}_2)^2 + (\mathcal{A}_3+\mathcal{A}_4)^2]^{1/2}$, and $\mathcal{D} = [(\mathcal{A}_1-\mathcal{A}_2)^2 + (\mathcal{A}_3 - \mathcal{A}_4)^2]^{1/2}$. With the application of external flux, the current and phase relationship [2, 3] for I_J^c may be found out. This will determine J_1^c and J_2^c and thereby δ and ξ. These determinations will predict the value of I_J^s. The spin current may be tested if the proposed set-up acts as a source of spin current.

The spin-Josephson current (9) is proportional to $\sin\xi$, which means that I_J^s is maximum when the orientation of the axes of non-centrosymmetry between the left and right side of the superconductors is transverse, and it vanishes when the axes are parallel. This particular property should be useful to control the spin-Josephson current by orienting the axes about which the inversion symmetries are lost. Therefore it will be useful to build a *Josephson spin switch*.

To summarize, we have shown that the Josephson tunnelling process between two superconductors without inversion symmetry consists of tunnelling of both spin-singlet and spin-triplet Cooper pairs. Because the direction-dependent tunnelling matrix element in a real junction and the momentum-dependent phases of the triplet components of the gap function, namely $\Delta_{\mathbf{k},\uparrow\uparrow}$ and $\Delta_{\mathbf{k},\downarrow\downarrow}$, are opposite, there is a net spin-Josephson current. This spin current is proportional to sine of the angle ξ between the axes of non-centrosymmetry in the bulk superconductors in each side of the junction. This may be used to make a *Josephson spin switch* with the controlling parameter ξ.

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