A superconductive magnetoresistive memory element using controlled exchange interaction

Sangjun Oh, D. Youm, and M. R. Beasley

Citation: Appl. Phys. Lett. **71**, 2376 (1997); doi: 10.1063/1.120032
View online: http://dx.doi.org/10.1063/1.120032
View Table of Contents: http://apl.aip.org/resource/1/APPLAB/v71/i16
Published by the American Institute of Physics.

A superconductive magnetoresistive memory element using controlled exchange interaction

Sangjun Oh and D. Youm
Department of Physics, Korea Advanced Institute of Science and Technology, Kusung-Dong, Yusung-Gu, Taejon 305-701, Korea

M. R. Beasley[a)]
Department of Applied Physics, Stanford University, Via Palou, Stanford, California 94305-4085

(Received 14 January 1997; accepted for publication 24 July 1997)

A memory device that can be switched between the normal state and superconducting state by an external magnetic field is proposed. The device consists of a superconducting/double magnetic (SM_1M_2) trilayer and is switched in a manner analogous to giant magnetoresistive memory devices. Using Usadel equations it is shown that the superconducting transition temperature of the device changes when the magnetic configurations of magnetizations of the two lower layers are switched between parallel and antiparallel. Appropriate design parameters are discussed and the materials issues analyzed. © *1997 American Institute of Physics.* [S0003-6951(97)00238-6]

Josephson junction logic circuits are well known for their high speed and extremely low power consumption. But Josephson junction memory circuits based on fluxoid quantization have not been so successful. Only relatively small memories have been fabricated.[1] Therefore, alternative memory concepts are of interest, particularly those with the potential to be fabricated in large sizes.

In this letter, we propose a superconducting memory device based on the superconducting/magnetic metal (SM) proximity effect. The structure of the device, shown in Fig. 1, is similar to that of giant magnetoresistance (GMR) memory devices and is switched in the same way as well.[2] It consists of a superconducting/double magnetic trilayer (SM_1M_2) structure in which the magnetization of the M_1 layer can be switched relative to the M_2 layer by the application of a small magnetic field. The physics of the device is completely different from GMR devices however. It is based on the oscillatory decay of the pair wave function predicted to occur in the M layer of a SM bilayer due to the influence of exchange interaction on the Cooper pairs.[3] Reversal of the magnetization in M_1 changes details of the decaying oscillations in the combined magnetic layers and thereby changes the transition temperature of the superconducting layer through the action of the proximity effect. Hence if the operating temperature is chosen appropriately, switching from the normal to the superconducting state can be achieved.

Thought of as a magnetoresistive switching element, this SM_1M_2 proximity device has 100% magnetoresistance (normal to superconducting transition) and lends itself to memory architectures similar to other magnetoresistive memories.

The theoretical operation of this device can be described using the Usadel equations with the inclusion of the exchange interaction. In our treatment, we include the important effects of spin-orbit scattering on the magnetic proximity effect recently clarified by Demler, Arnold, and Beasley.[4] We assume the dirty limit both for simplicity and because it is the likely situation in practice. Following Ref. 4, we take

$$\frac{\hbar}{2} D_s \frac{d^2}{dx^2} f^s_\pm(\omega,x) - \hbar|\omega| f^s_\pm + \Delta = 0\,,$$
$$\Delta(x) = \frac{\lambda}{2}\sum_\omega [f^s_+(\omega,x) + f^s_-(\omega,x)]\,, \tag{1}$$

where f^s_+, f^s_- are the anomalous Green functions for the spin up and the spin down states, D_s is the diffusion coefficient, and Δ is the superconducting order parameter in the superconducting layer, and

$$\frac{\hbar}{2} D_m \frac{d^2}{dx^2} f^m_\pm(\omega,x) \mp sgn(\omega) ih f^m_\pm = \frac{1}{\tau_{so}}(f^m_\pm - f^m_\mp)\,, \tag{2}$$

where $f^m_\pm$ are the anomalous Green functions, D_m is the diffusion constant, h is the exchange energy, and τ_{so} is the spin-orbit relaxation time in the ferromagnetic layers. λ is the superconducting coupling constant, λ_0, times $\pi k_B T$. We assume λ is zero in the magnetic layers. At the two boundaries, $x=0$ and $x=$d (see Fig. 1), f and $\sigma(d/dx)f$ must be

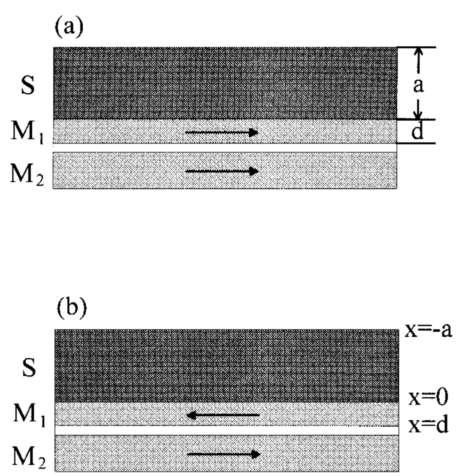

FIG. 1. Structure of the superconducting/double magnetic trilayer (SM_1M_2). (a) Parallel and (b) antiparallel configurations of the magnetization of M_1 and M_2 layers. In the calculations, we assume the thickness of M_2 layer to be semi-infinite and only the thickness, d, of the M_1 layer is varied.

[a)]Electronic mail: beasley@ee.stanford.edu

continuous, where σ is the normal state conductivity of the appropriate layer. Note that in Eqs. (1) and (2) we neglect the conventional normal metal proximity effect in the M layers. If we introduce the more convenient functions,

$$F_+ = f_+ + f_-, \quad F_- = f_+ - f_-, \tag{3}$$

Eqs. (1) and (2) become

$$\frac{\hbar}{2} D_s \frac{d^2}{dx^2} F_+^s(\omega,x) - \hbar|\omega| F_+^s + 2\Delta = 0,$$

$$\frac{\hbar}{2} D_s \frac{d^2}{dx^2} F_-^s(\omega,x) - \hbar|\omega| F_-^s = 0,$$

$$\frac{\hbar}{2} D_m \frac{d^2}{dx^2} F_+^m(\omega,x) - sgn(\omega) ih F_-^m = 0, \tag{4}$$

$$\frac{\hbar}{2} D_m \frac{d^2}{dx^2} F_-^m(\omega,x) - sgn(\omega) ih F_+^m = \frac{2}{\tau_{so}} F_-^m,$$

$$\Delta(x) = \frac{\lambda}{2} \sum_\omega F_+^s.$$

Assuming M_2 is effectively semi-infinite, we have $\lim_{x\to\infty} F_\pm^{m_2}(x) = 0$. We now look for a solution of the form:

$$F_-^{m_1} = A_1 e^{-k_m x} + A_2 e^{-k_m^* x} + A_3 e^{k_m x} + A_4 e^{k_m^* x},$$

$$F_+^{m_1} = -\alpha A_1 e^{-k_m x} + \alpha^* A_2 e^{-k_m^* x} - \alpha A_3 e^{k_m x} + \alpha^* A_4 e^{k_m^* x}, \tag{5}$$

$$F_-^{m_2} = C_1 e^{-k_m x} + C_2 e^{-k_m^* x},$$

$$F_+^{m_2} = \alpha C_1 e^{-k_m x} - \alpha^* C_2 e^{-k_m^* x},$$

and

$$F_-^s = C \cosh k_s(x+a), \tag{6}$$

$$F_+^s = \Delta(x) f(\omega),$$

$$\Delta(x) = U \cos \beta(x+a), \quad f(\omega) = \frac{4}{\hbar(2|\omega| + D_s \beta^2)}, \tag{7}$$

for magnetic layers M_1, M_2 and for the superconducting layer, respectively. C_1, C_2, A_1, A_2, A_3, A_4, and C, U are numerical coefficients and

$$k_s = \sqrt{\frac{2|\omega|}{D_s}}, \quad k_m = \frac{1}{\sqrt{\hbar D_m}} \left(\sqrt{h + \frac{1}{\tau_{so}}} + i \sqrt{h - \frac{1}{\tau_{so}}} \right),$$

$$\alpha = \frac{sgn(\omega) ih}{\frac{\hbar}{2} D_m k_m^2}.$$

The assumed factorized form of $F_+^s(x,\omega)$ in Eq. (6) also requires that β be independent of ω. The conditions for the validity of this condition are discussed below.

Taking for simplicity $\sigma_{m_1} = \sigma_{m_2} = \sigma_m$, application of the boundary condition at $x = d$ yields the following relations:

$$\begin{aligned} \frac{\sigma_s}{\sigma_m} k_s \tanh(k_s a) &= \frac{[k_m(1+R) + k_m^* S^*]X + [k_m^*(1+R^*) + k_m S]}{(1-R-S^*)X + (1-R^*-S)}, \\ -\frac{\sigma_s}{\sigma_m} \beta \tan(\beta a) &= \frac{-[\alpha k_m(1+R) - \alpha^* k_m^* S^*]X + [\alpha^* k_m^*(1+R^*) - \alpha k_m S]}{-[\alpha(1-R) + \alpha^* S^*]X + [\alpha^*(1-R^*) + \alpha S]}, \end{aligned} \tag{8}$$

where we define R, S as $R = (k_m^* + k_m)/(k_m^* - k_m) e^{2k_m d}$, $S = -[(\alpha^* - \alpha) k_m^* / (k_m^* - k_m)\alpha] e^{(k_m + k_m^*)d}$ for convenience, and where $X = A_3/A_4$ can be eliminated after combining the above two equations. Then we can express β in terms of the given parameters, a, d, h, D_m, D_s, σ_s, σ_m, and τ_{so}. We also define the coherence lengths

$$\xi_m = \sqrt{\frac{4\hbar D_m}{h}}, \quad \xi_s = \sqrt{\frac{\hbar D_s}{2\pi k_B T_{co}}}.$$

Detailed examination of the theory shows that Eq. (6) is only valid in the limits, $\epsilon = (\sigma_s/\sigma_m)(\xi_m/\xi_s) \gg 1$, or $\epsilon \ll 1$. Note that since $\sigma = eN(0)D$, ϵ may also be written as

$$\epsilon = 2 \frac{N_s(0)}{N_m(0)} \sqrt{\frac{D_s}{D_m}} \sqrt{\frac{2\pi k T_{co}}{h}}.$$

Thus our calculations are similar to those of Buzdin and Kupriyanov[3] and are meaningful only in the limiting cases indicated. For general values of ϵ a numerical solution is required.[5] The dimensionless quantity $\beta\xi_s$ can be expressed neatly in terms of the dimensionless quantities, ϵ, d/ξ_m, a/ξ_s, $h\tau_{so}$. Combined with $1 = \pi k T_{co} \lambda_o \Sigma_\omega (1/\hbar|\omega|)$ and Eq. (9), Eq. (4) gives

$$\ln t = \Psi\left(\frac{1}{2}\right) - \Psi\left(\frac{1}{2} + \frac{\beta^2 \xi_s^2}{2t}\right), \tag{9}$$

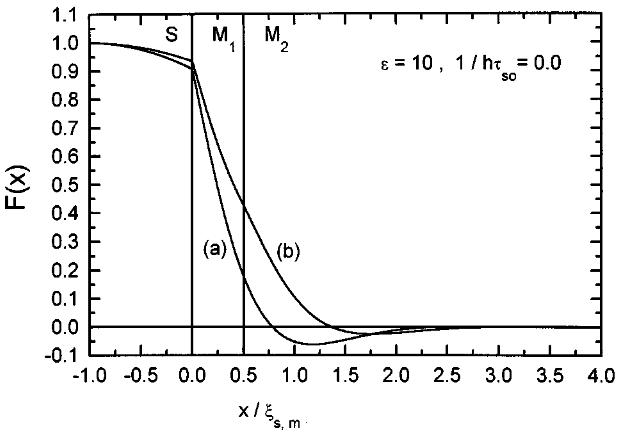

FIG. 2. Normalized pair wave functions, $F(x)=\Sigma_\omega[f_-(\omega,x)+f_+(\omega,x)]$, as functions of the normalized spatial coordinate $x/\xi_{s,m}$. Curves (a) and (b) are for cases (a) and (b) in Fig. 1 where the thickness of the S layer, a, is ξ_s, and the thickness of the M_1 layer, d, is $\frac{1}{2}\xi_m$.

where $t=T_c/T_{co}$, T_{co} is the critical temperature of the bulk material and Ψ is the digamma function.

The effect of exchange field on the spatial dependence of the pair wave function $F(x)$ is shown in Fig. 2, for $\epsilon=10$ and $1/h\tau_{so}=0$. Curves (a) and (b) are for parallel and antiparallel configurations of M_1 and M_2, respectively. From the previous expression of the wave vector k_m, the increase of the wave length of oscillation in the curve (b) implies a decrease of the effective exchange energy, h, which can be easily understood. In the antiparallel configuration the different signs of exchange energies in the M_1 and M_2 layers make its average value small. This feature can be interpreted as a reduction of the pair breaking effect.

Figure 3 shows a plot of t vs d/ξ_m at a fixed $a/\xi_s=0.7$ for various spin-orbit scattering rates $1/h\tau_{so}$. As expected, spin-orbit scattering reduces the oscillatory part pair breaking effect, hence increasing $T_c(min)$ and reducing the peak in T_c vs d/ξ_m. Spin-orbit scattering, in effect, mixes the spin-up and spin-down states of the individual electrons in a Cooper pair, thereby negating any effect of the exchange field. A similar effect is well known in the case of Pauli limiting in the upper critical field of a superconductor.[6]

Practical application favors large ϵ. From a materials point of view this implies the need for a relatively clean superconductor and a dirty magnetic metal, consistent with small spin-orbit scattering. This in turn implies the need for magnetic metals incorporating only elements with small atomic numbers. The requirement for large ϵ also favors magnetic metals with low exchange energies. Low exchange energies have the additional advantage that ξ_m increases, which means thicker films can be used and this is consistent with requirement that $d/\xi_m=0.4$.

As previously noted, the physical basis of this device is closely related to that associated with T_c oscillations in SM bilayers.[5] Attempts to experimentally observe this effect using elemental magnetic materials [Fe,[7] Ni,[8] and Gd (Ref. 9)] are controversial and inconclusive, despite favorable estimates of ϵ (5.1 for V/Fe SM bilayers and 3.7 for Gd/Nd). On the other hand, the theory appears to be on a firm basis. We speculate that the problem may be the need for magnetic materials with larger ξ_m for purely practical reasons. ξ_m of Fe is less than 6 Å,[7] and Ni films thinner than 8 Å are nonmagnetic.[8] Also, Gd may suffer from large spin-orbit scattering. The outlook may be brighter with magnetic alloys or compounds with low exchange fields (small Curie temperatures) formed from small Z elements.

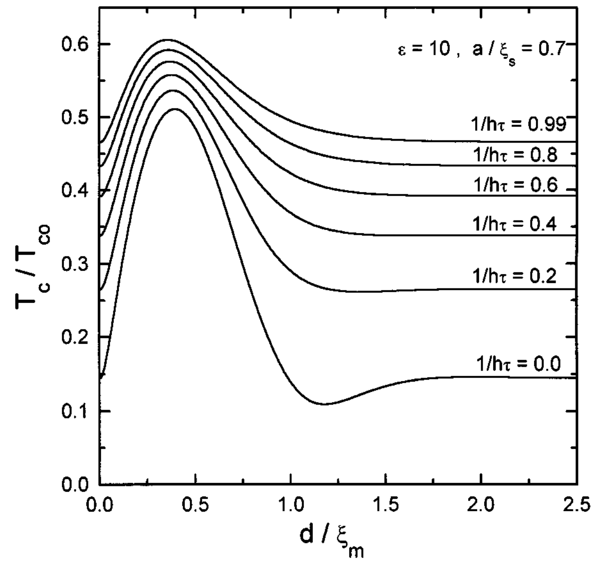

FIG. 3. T_c/T_{co} vs d/ξ_m at a fixed $a/\xi_s=0.7$ and $\epsilon=10$ for various spin-orbit scattering rates $1/h\tau_{so}$. For $\epsilon<1$, the effect is small.

The work at Stanford was supported by the DOD URI program in Superconductivity Electronics and by KMOST.

[1] M. G. Forrester, J. X. Przybysz, J. Talvacchio, J. Kang, A. Davidson, and J. R. Gavaler, IEEE Trans. Appl. Supercond. **5**, 3401 (1995).
[2] J. M. Daughton, Thin Solid Films **216**, 162 (1992); K. T. M. Ranmuthu,I. W. Ranmuth, A. V. Pohm, C. S. Comstock, and M. Hassoun, IEEE Trans. Magn. **28**, 2359 (1992); A. V. Pohm, C. S. Comstock, and J. M. Daughton, IEEE Trans. Magn. **25**, 4266 (1989).
[3] A. I. Buzdin and M. Yu. Kupriyanov, JETP Lett. **52**, 487 (1990); **53**, 321 (1991).
[4] E. A. Demler, G. B. Arnold, and M. R. Beasley, Phys. Rev. B **55**, 15174 (1997).
[5] Z. Radovic, M. Ledvij, L. Dobrosavljevic-Grujic, A. I. Buzdin, and J. R. Clem, Phys. Rev. B **44**, 759 (1991).
[6] H. Engler and P. Fulde, Phys. Kondens. Mater. **7**, 150 (1968).
[7] P. Koorevaar, Y. Suzki, R. Coehoorn, and J. Aarts, Phys. Rev. B **49**, 441 (1994);Th. Muhge, N. N. Garifyanov, Yu. V. Goryunov, G. G. Khaliullin, L. R. Tagirov, K. Westerholt, I. A. Garifullin, and H. Zabel, Phys. Rev. Lett. **77**, 1857 (1996).
[8] H. Homma, C. S. L. Chun, G. G. Zheng, and I. K. Schuller, Phys. Rev. B **33**, 3562 (1986).
[9] C. Strunk, C. Surgers, U. Paschen, and H. v. Lohneysen, Phys. Rev. B **49**, 4053 (1994); J. S. Jiang, D. Davidovic, Daniel H. Reich, and C. L. Chien, Phys. Rev. Lett. **74**, 314 (1995).