Magnetoelectric effect in helical metallic antiferromagnets

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The magnitude of the kinetic magnetoelectric effect in an incommensurate structure is found. The Fröhlich conductivity mechanism for the antiferromagnetic phase of certain rare-earth elements and compounds is discussed.

In several antiferromagnetic insulators the magnetoelectric effect (the appearance of magnetization proportional to the electric field) can be produced by disrupting the invariance with respect to the sign of the time and the spatial inversion.[1,2] The helical structure in this sense has the necessary symmetry properties.

We will show, however, that the magnetoelectric effect in helical metallic antiferromagnets is determined by the kinetic phenomena. In other words, in the case of conductors with stereoisomerism, the key factor in this phenomenon is, as in Ref. 3, the disruption of the invariance with respect to the time reversal due to the dissipative nature of the current flow.

The helical phase has been established in several rare-earth elements and their compounds (in Dy, Ho, and Eu, for example) by means of neutron-diffraction analysis. Furthermore, the antiferromagnetic vector **Q** of these compounds is essentially incommensurate with the lattice vector. The origin of the incommensurate structure is usually attributed (see, e.g., Refs. 4–6) to the presence on the electronic Fermi surface of congruent regions which are linked by the Keldysh-Kopaev relation

$$\epsilon(\mathbf{p} + \mathbf{Q}) = -\epsilon(\mathbf{p}) \tag{1}$$

(these regions are shown schematically in Fig. 1). Their combination due to the formation of a superstructure should lead to the development of energy gaps on a part of the Fermi surface. Evidence for this conclusion comes from the fact that the resistance of Dy[7] and Ho[8], for example, increases appreciably below the Néel point.

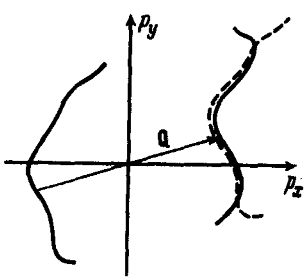

FIG. 1. Superimposed parts of the Fermi surface.

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In systems with local moments, the Ruderman-Kittel-Kasuya-Yosida interaction accounts for the antiferromagnetic order and in systems without local moments the formation of an electron spin-density wave accounts for the antiferromagnetic order. In each case the transition mechanism is attributed, as has been suggested, to the instability of spectrum (1) with respect to the gap formation. Below the transition point, components of the nondiagonal electron order parameter, $\hat{\Delta}_{RL}$ and $\hat{\Delta}_{LR}$, of the following type arise:

$$\hat{\Delta}_{RL} = (\vec{\sigma}\mathbf{d})\exp(i\mathbf{Qr}), \tag{2}$$

These components are proportional to $(\hat{\Delta}_{RL})_{\alpha\beta} \sim \langle \psi_{R\alpha}\psi^{+}_{L\beta}\rangle$ and $(\hat{\Delta}_{LR})_{\alpha\beta} \sim \langle \psi_{L\alpha}\psi^{+}_{R\beta}\rangle$, respectively (the subscripts R and L correspond to the right and left parts of the Fermi surface in Fig. 1). Accordingly,

$$(\hat{\Delta}_{RL})^{+} = \hat{\Delta}_{LR}\ . \tag{3}$$

The splitting of the energy spectrum (ϵ) is determined by the matrix determinant of the form

$$\begin{pmatrix} \xi - \epsilon & \hat{\Delta}_{RL} \\ \hat{\Delta}^{+}_{RL} & -\xi - \epsilon \end{pmatrix},$$

where $\xi = v_F(p - p_F)$ is the energy separation from the Fermi surface. Evaluating the determinant of the right part of the Fermi surface, we find

$$\epsilon^2_R = \xi^2 + (\mathbf{d}\mathbf{d}^*) + i(\vec{\sigma}[\mathbf{d}\times\mathbf{d}^*])\ . \tag{4}$$

If the vector $\mathbf{d}$ is real ($\mathbf{d}^2 \neq 0$), the structure corresponds to a sinusoidal wave. The spectrum will then have, according to (4), a gap for each spin component.

A purely helical wave corresponds to $\mathbf{d}^2 = 0$: $\mathbf{d} = \mathbf{d}_1 + i\mathbf{d}_2$, where $\mathbf{d}_1 \perp \mathbf{d}_2$ and $|\mathbf{d}_1| = |\mathbf{d}_2|$. In this case the spectrum has one branch with a gap and one branch without it:

$$\epsilon^2_R = \xi^2 + |\mathbf{d}|^2(1 + \sigma_z), \tag{5}$$

where the direction z is along the helicon axis:

$$\mathbf{n} \parallel [\mathbf{d}_1 \times \mathbf{d}_2]\ .$$

For the left part of the Fermi surface we find

$$\epsilon^2_L = \xi^2 + |\mathbf{d}|^2(1 - \sigma_z)\ . \tag{5'}$$

Spectra (5) and (5') for different spin orientations relative to the vector $\mathbf{n}$ [we assume that spectrum (1) corresponds to a common energy branch] are illustrated in Fig. 2. At low temperatures we can thus clearly see the origin of the effect.

In the presence of an electric current the occupation numbers are redistributed, as indicated by the heavy line in Fig. 2. The redistribution, however, affects only the gap-free branches. To the right and left, these branches have different spin directions, giving rise to magnetization

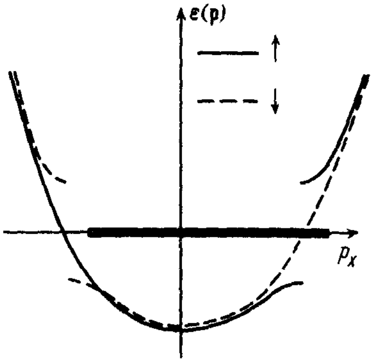

FIG. 2. Spectrum of electronic excitations in the helical phase. The heavy line shows the redistribution of the occupation numbers in the presence of an electric current.

$$M_i = \alpha_{ik} E_k \, . \tag{6}$$

In the geometry in Fig. 1, **M** is directed along **n**:

$$\mathbf{M} = \alpha \mathbf{e}_z E_x \tag{6'}$$

(in the exchange approximation the helicon axis is not specified).

We now give the results of microscopic calculations. At $T = 0(|\mathbf{d}| \sim T_N)$ for α we find

$$\alpha = -\mu_B e S \frac{2\sqrt{2}|\mathbf{d}|\tau^2_{imp}}{\pi^3} \, . \tag{7}$$

In (7) τ_{imp} is the transport time (the scattering time with a transfer of momentum $\sim \mathbf{Q}$) due to the impurities, and S is the area of each region of the Fermi surface. At higher temperatures the dissipative mechanism stems from the inelastic (phonon) processes. In the neighborhood of T_N we find

$$\alpha = -\mu_B e S \frac{(|\mathbf{d}|\tau)^2}{(2\pi)^2 T_N (1 + 8(|\mathbf{d}|\tau)^2)^{1/2}}$$

$$|\mathbf{d}| \ll T_N \, . \tag{8}$$

(For impurities such an expression predicts a spasmodic increase in α below T_N, which stems from the absence of scattering mechanisms which mix up the spin channels.) The spin structure of the antiferromagnets under discussion has an essentially incommensurate period. The wave can "slip" along the vector **Q**, transferring the charge. The spin structure is, of course, pinned by the impurities. In the exchange approximation, the pinning may be, in our view, not too strong (it is attributable to the Gaussian fluctuations in a uniform distribution of impurities, since the impurities shift the Néel temperature). That the theoretical understanding of the nature of the superstructure which we discussed at the beginning of this paper is correct can be verified experimentally by measuring the conductivity in strong electric fields or in a microwave field. Leaving the explicit expressions for the conductivity under conditions of a pinned and a traveling wave to be written in a detailed paper, we note that the results are generally similar to the situation observed in transition-metal trichalcogenides (see Refs. 10 and 11, for example). The entrainment of spin-density waves by an electric field in Dy or Ho, for example, should lead to a smoothing out of the resistance-induced anomalies near T_N, which were observed in Refs. 7 and 8, and to the appearance of a nonlinear current-voltage characteristics near the threshold electric fields.

Estimate of the expression for the magnetoelectric coefficient near T_N when the "helix" is moving along the vector **Q** is

$$\alpha \sim \mu_B eS \frac{|\mathbf{d}|^2 \tau}{T_N^2} .$$

Comparing this result with (8), we find that the field-induced deflection of the spin-density wave reduces the magnetoelectric effect appreciably and even changes its sign.

In general, magnetic moment (6) is small in comparison with the magnetic field of the current flow. In light of the discussion above, it would be of considerable interest to study the above-mentioned nonlinear and frequency effects in the conductivity.

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