Circular dichroism in high-temperature superconductors: A characteristic of broken space-inversion symmetry

Badredine Arfi
Department of Physics, University of Illinois at Urbana–Champaign, and Materials Research Laboratory, 1110 West Green Street, Urbana, Illinois 61801

L. P. Gor'kov
Department of Physics, University of Illinois at Urbana–Champaign, 1110 West Green Street, Urbana, Illinois 61801 and L. D. Landau Institute for Theoretical Physics, Chernogolovka, Moscow District, 142434, Russia

(Received 2 March 1992)

The circular optical dichroism recently observed in a number of high-temperature superconductors has been commonly interpreted in terms of violation of both time-reversal invariance and space-inversion symmetry (the magnetoelectric effect). It is shown, however, that in view of numerous experimental observations of various piezoelectric or ferroelectric phenomena in these superconductors, the effect could be naturally interpreted in terms of a linear space dispersion of the dielectric tensor of a dissipative medium that lacks a center of inversion. Both bulk and surface contributions to the effect are discussed.

I. INTRODUCTION

The recent interest in investigating the possible occurrence of circular optical dichroism in certain high-temperature superconductors[1] was initiated by a remark made by Wen and Zee,[2] namely, that the magnetoelectric effect,[3] which would be, thermodynamically speaking, a consequence of broken time-reversal and space-inversion symmetries, could be inherent to models of the so-called "anyon" superconductivity.[4]

Among high-T_c superconductors circular dichroism was first observed in $YBa_2Cu_3O_7$ compound.[1] A comparable and even bigger effect was also reported for the cubic compound $Ba_{1-x}Rb_xBiO_3$ in the normal state.[1] The CD rotation angle arises slightly above 200 K and has a pronounced onset and temperature dependence. In addition to this, a large amount of literature[5–9] reports on the observation of various piezoelectric, pyroelectric, or ferroelectric-related phenomena in both films and single crystals of $YBa_2Cu_3O_7$ and La_2CuO_4 compounds. In spite of the numerous precautions taken in the experiments, it is still unclear to what extent in *each* experiment, the role played by the sample surfaces is important. These experiments, if taken altogether, seem to indicate that as the temperature is decreased from room temperature most superconducting perovskites undergo a transition to a phase lacking a center of inversion. For example, in the case of $Ba_{1-x}Rb_xBiO_3$ compound a transition to the new "phase," as the CD experiments seem to imply,[1] would have an onset temperature of ~ 200 K. If this point of view is accepted, the magnetoelectric effect (and consequently the circular dichroism) could then be understood as being of a kinetic origin, as explained below.

It was emphasized in Refs. 10 and 11 that for the magnetoelectric effect to exist in thermodynamical equilibrium (i.e., with both time-reversal and parity symmetries broken), the magnetic cell is to coincide with the crystalline one. This restriction is not necessary if the phenomenon has a kinetic origin, i.e., if dissipation substitutes for the breaking of time-reversal symmetry in the system. An example of magnetoelectric effect in conductors having helical magnetic structure has been considered by one of the authors.[12] Another realization of kinetic magnetoelectric effect has been considered in Ref. 13 in the case of a conductor possessing mirror isomer symmetry. These authors used a model put forward by Belinicher and Sturman[14] in which a conductor lacking a center of inversion is simulated by an ordered arrangement of impurities whose scattering potential is asymmetric. Coming back to the problem of circular dichroism in the high-temperature superconductors, as it is posed in the abstract, we shall see below that after a minor reformulation it is possible to accommodate the results of Ref. 13 and 14 to our needs.

A few words are needed to establish a correspondence between the above properties (the hypothetical noncentrosymmetric phase) in the new superconductors in the normal state and the model of Refs. 13 and 14. The actual source of resistivity ($\rho \sim T$) in the new materials is still a controversial issue. One can only guess that, since $1/\tau \sim T$, this is somehow due to some sort of quasielastic scattering and can be considered by the same methods as for the impurity scattering. In addition to that, layered cuprates (La_2CuO_4 or $YBa_2Cu_3O_{7-\delta}$ have defects of structural type, especially near the sample surface. For $Ba_{1-x}Rb_xBiO_3$ substitutional defects are inherent in the range of $x \sim 0.4$. A trivial comment of a rather general character is that, rigorously speaking, there are no reasons for a local defect to possess the local symmetry of the host lattice. The overall symmetry of the crystal containing substitutional impurities (say for example in

$Ba_{1-x}Rb_xBiO_3$) can remain cubic, since averaging means not only space homogeneity but also averaging over all configurations of defect shapes allowed by cubic lattice symmetry. If a structural transition takes place in the lattice, it, correspondingly, produces a preferential orientation of the scattering centers. At higher defect concentrations, defect ordering is possible.

Most experiments in which attempts were made to observe the aforementioned properties, usually use various pulse techniques or high-frequency illumination. In the remaining part of the paper, we consider much lower frequencies. We shall first study the magnetoelectric effect and, hence, the CD in the normal skin regime ($\omega \ll 1/\tau$) in which case the *bulk* properties would be measured. In addition, we shall also consider the regime of anomalous skin effect where both bulk and surface sources of dissipation are present on equal footing. In the so-called far-infrared regime, the only source of dissipation is scattering on the *surface roughness*. The paper is organized as follows. In Sec. II, we discuss the kinetic magnetoelectric effect as well as the photogalvanic effect, while in Sec. III we calculate the circular optical dichroism rotation angle in three different cases: a first case where the effect has a bulk origin, a second case that corresponds to the anomalous skin effect regime, and finally a third case that corresponds to the infrared regime and where the surface roughness is assumed to play a dominant role in the scattering process. In Sec. IV, we draw our conclusions.

II. KINETIC MAGNETOELECTRIC EFFECT AND PHOTOGALVANIC EFFECT

The magnetoelectric effect, known already for years,[10] is that an electric field **E** induces a macroscopic magnetic moment **M** when the above-mentioned symmetries are broken:

$$M_i = \alpha_{ij} E_j \ . \tag{1}$$

In general α is a second-rank tensor. In the cubic symmetry case, α reduces to a pseudoscalar.

In the case of conducting materials, the applied electric field induces currents and any experimental observation of relation (1) requires an AC regime. The macroscopic electrodynamic equations describing the medium are[3]

$$\nabla\times\mathbf{E} = -\frac{1}{c}\frac{\partial \mathbf{B}}{\partial t}, \quad \nabla\cdot\mathbf{B}=0 \ , \tag{2}$$

$$\nabla\times\mathbf{B} = \frac{1}{c}\frac{\partial \mathbf{D}}{\partial t}, \quad \nabla\cdot\mathbf{D}=0 \ . \tag{2'}$$

In writing down the general form of the dielectric constant, we take into account the linear space dispersion only. This translates in Fourier space into a linear **k** dependence, namely

$$\epsilon_{ij}(\omega,\mathbf{k}) = \epsilon_{ij}^{(0)}(\omega) + i\gamma_{ijl}k_l \ , \tag{3}$$

where ω is the frequency of the applied electric field. Let us now briefly generalize the symmetry relations obeyed by ϵ_{ij} in transparent media to the case where dissipation is present. The first relation

$$\epsilon_{ij}(-\omega,-\mathbf{k}) = \epsilon_{ij}^{*}(\omega,\mathbf{k}) \tag{4}$$

follows from the definition $\mathbf{D}=\epsilon\mathbf{E}$, in which the dielectric tensor establishes a linear relation between the two real (in space) fields, **D** and **E**. The second relation

$$\epsilon_{ij}(\omega,\mathbf{k}) = \epsilon_{ji}(\omega,-\mathbf{k}) \tag{5}$$

is a consequence of Onsager's symmetry principle for generalized kinetic coefficients. In the absence of dissipation, γ_{ijl} would be real.[3] Instead of Eq. (3), we rewrite $\epsilon_{ij}(\omega,\mathbf{k})$ in the following more standard form:

$$\epsilon_{ij}(\omega,\mathbf{k}) = \epsilon_{ij}(\omega) - \frac{4\pi}{\omega}\sigma_{ijl}k_l \tag{6}$$

and investigate the symmetry of $\sigma_{ijl}(\omega)$ as it would stem from Eqs. (4) and (5). In Eq. (6), $\epsilon_{ij}(\omega)$ includes the usual wave-vector-independent contribution to the conductivity. The real part of γ_{ijl} is omitted, since it is irrelevant to the following considerations. Combining Eqs. (4) and (5) one arrives at

$$\sigma_{ijl}(\omega) = \sigma_{ijl}^{*}(-\omega), \quad \sigma_{ijl}(\omega) = -\sigma_{ilj}(\omega) \ .$$

In the limit $\omega \to 0$, it follows from the above relations that

$$\sigma_{ijl} = i\chi_{ijl}, \quad \chi_{ijl} = -\chi_{jil} \ . \tag{7}$$

(It is appropriate to remember here that a nonzero third-rank tensor can only exist for media without a center of symmetry). Introducing the unit vector $\mathbf{n}=\mathbf{k}/|\mathbf{k}|$, it is straightforward to express the asymmetric tensor $-(4\pi/\omega)\sigma_{ijl}n_l|\mathbf{k}|$ in terms of the axial vector g_{ml},

$$-\frac{4\pi}{\omega}ik_l\chi_{ijl} = -\frac{4\pi}{\omega}i|\mathbf{k}|\epsilon_{ijm}g_{ml}n_l \ .$$

We can write

$$g_{xx} = \chi_{yzx}, \quad g_{xy} = \chi_{yzy}, \quad g_{yx} = \chi_{zxx} \ \text{etc} \ .$$

In the case of the cubic symmetry, g_{ml} reduces to a pseudoscalar $g_{ml} = d\delta_{ml}$. Using Eqs. (1) and (2′) we obtain

$$\mathbf{M} = \alpha\mathbf{E} \ . \tag{1'}$$

With all this in mind, we can, after minor modifications, utilize the results of Ref. 13. It is shown there that the major contribution to the current that would result in Eq. (1′), is due to the central cross section of the diagram shown in Fig. 1. The explicit form of the coefficient in Eq. (1) in the case, for instance, of cubic symmetry is

$$\alpha = \frac{1}{12\pi^2}\frac{e^2}{\hbar c}(p_F l)^2\left(\tau n_i \frac{m p_F}{\pi}\right)^2 \times \int \frac{dO_p dO_{p'} dO_{p''}}{(4\pi)^3}\hat{\mathbf{p}}\cdot(\hat{\mathbf{p}}'\times\hat{\mathbf{p}}'')W_{\mathbf{pp}'}W^{*}_{\mathbf{p}'\mathbf{p}}W_{\mathbf{p}'\mathbf{p}''}W^{*}_{\mathbf{p}''\mathbf{p}'} \ , \tag{8}$$

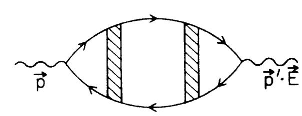

FIG. 1. Diagram giving leading contribution to the coefficient α. The hatched bars are the scattering amplitudes.

where $W_{\mathbf{pp}'}$ is the amplitude for the electron to scatter from the state characterized by $\mathbf{p}$ to the one characterized by $\mathbf{p}'$, and τ is the scattering relaxation time given by

$$\frac{1}{\tau}=n_i\frac{mp_F}{\pi}\int\frac{dO_p}{4\pi}|W_{\mathbf{pp}}|^2 .$$

The symmetry of the scattering potential can be seen through the behavior of, say, $W_{\mathbf{pp}'}$ under the symmetry operations of the cubic group (remember that the symmetry transformations are applied simultaneously to $\mathbf{p}$ and $\mathbf{p}'$ vector directions on the Fermi surface). Let i be the specific representation for the order parameter in the phase of lower symmetry. Since the ordering of the impurities is induced by a lowering of the cubic symmetry, one can write in a general way

$$W_{\mathbf{pp}'}=W^s_{\mathbf{pp}'}+W^i_{\mathbf{pp}'} . \quad (9)$$

In this equation $W^s_{\mathbf{pp}'}$ is the invariant (symmetric) part, while $W^i_{\mathbf{p}'\mathbf{p}}$ transforms in the same way as the order parameter under rotations and is asymmetric ($W^i_{\mathbf{pp}'}\neq W^i_{\mathbf{p}'\mathbf{p}}$). Having in mind the numerous experimental observations of various piezoelectric and/or ferroelectric phenomena in the high-temperature superconductors, which is a signature of the absence of a center of inversion, we take $W^i_{\mathbf{pp}'}$ to be odd under space inversion. Therefore, we can write the following relations

$$W^s_{\mathbf{pp}'}=W^s_{\mathbf{p}'\mathbf{p}}=-W^s_{-\mathbf{p}-\mathbf{p}'} ,$$

$$W^i_{\mathbf{pp}'}=-W^i_{\mathbf{p}'\mathbf{p}}=-W^i_{-\mathbf{p}-\mathbf{p}'} . \quad (9')$$

Because $\hat{\mathbf{p}}\cdot(\hat{\mathbf{p}}'\times\hat{\mathbf{p}}'')$ is a pseudoscalar, the nonvanishing contribution to α [using Eqs. (9′)], is found to be

$$\alpha=\frac{1}{12\pi^2}\frac{e^2}{\hbar c}(p_Fl)^2\left[\tau n_i\frac{mp_F}{\pi}\right]^2\frac{2}{3}\int\frac{dO_pdO_{p'}dO_{p''}}{(4\pi)^3}\hat{\mathbf{p}}\cdot(\hat{\mathbf{p}}'\times\hat{\mathbf{p}}'')i$$
$$\times\mathrm{Im}\{|W^i_{\mathbf{p}''\mathbf{p}'}|^2W^{s*}_{\mathbf{pp}'}W^i_{\mathbf{pp}'}-|W^i_{\mathbf{pp}'}|^2W^{s*}_{\mathbf{p}''\mathbf{p}'}W^i_{\mathbf{p}''\mathbf{p}'}+|W^i_{\mathbf{p}'\mathbf{p}}|^2W^{s*}_{\mathbf{p}''\mathbf{p}}W^i_{\mathbf{p}''\mathbf{p}}$$
$$-|W^i_{\mathbf{p}''\mathbf{p}}|^2W^{s*}_{\mathbf{p}'\mathbf{p}}W^i_{\mathbf{p}'\mathbf{p}}+|W^i_{\mathbf{pp}''}|^2W^{s*}_{\mathbf{p}'\mathbf{p}''}W^i_{\mathbf{p}'\mathbf{p}''}-|W^i_{\mathbf{p}'\mathbf{p}''}|^2W^{s*}_{\mathbf{pp}''}W^i_{\mathbf{pp}''}\} . \quad (10)$$

Written in this way, it is clear that the expression between curly brackets cancels the antisymmetry character of $\hat{\mathbf{p}}\cdot(\hat{\mathbf{p}}'\times\hat{\mathbf{p}}'')$.

Quite generally, one can expand both $W^s_{\mathbf{pp}'}$ and $W^i_{\mathbf{pp}'}$ in the complete set of some basis functions:

$$W^s_{\mathbf{pp}'}=\sum_{K,n}\psi^K_n(\mathbf{p})\psi^K_n(\mathbf{p}')A^K_s , \quad (11)$$

$$W^i_{\mathbf{pp}'}=\sum_{n,i=l\otimes m}[\psi^l_n(\mathbf{p})\psi^m_n(\mathbf{p}')-\psi^l_n(\mathbf{p}')\psi^m_n(\mathbf{p})]B^{lm}_i . \quad (11')$$

The upper index enumerates the specific representations of the cubic group, while n takes into account the fact that for the point group the complete basis set includes an infinite number of eigenfunctions transforming in accordance with each of the ten representations of O_h. In Eq. (11′) the summation over l and m is restricted by the only condition that the antisymmetrized product $[\psi^l,\psi^m]$ is to transform according to the chosen representation i of the order parameter (one of these is to be odd and the other one even with respect to inversion). In such a form, Eqs. (11) and (11′) contain both the symmetric and asymmetric parts of $W_{\mathbf{pp}'}$.[14]

In the case of the orthorhombic group D_{2h} (relevant to $YBa_2Cu_3O_7$ compound) the coefficient $\overleftrightarrow{\alpha}$ is a tensor whose components are given by

$$\alpha_{ii}=\frac{i}{2c}\epsilon_{jlm}\sigma_{jlm} ,$$
$$\alpha_{ij}+\alpha_{rm}\epsilon_{lmi}\epsilon_{lrj}=\frac{i}{c}\epsilon_{jrl}\sigma_{irl} . \quad (12)$$

Expressions (11) and (11′) are still valid with the modification appropriate to the eight representations of D_{2h}. We note that Eqs. (12) are of a general character and can be used to find the finite components of the tensor α for any symmetry group. In the case of the group D_{2h}, we rewrite Eqs. (11) and (11′) in the following form

$$W^s_{\mathbf{pp}'}=\sum_{K,n}[\psi^{(u)K}_n(\mathbf{p})\psi^{(u)K}_n(\mathbf{p}')+\psi^{(g)K}_n(\mathbf{p})\psi^{(g)K}_n(\mathbf{p}')]A^K_s , \quad (13)$$

$$W^i_{\mathbf{pp}'}=\sum_{n,i=l\otimes m}[\psi^{(u)l}_n(\mathbf{p})\psi^{(g)m}_n(\mathbf{p}')-\psi^{(g)l}_n(\mathbf{p})\psi^{(u)m}_n(\mathbf{p}')]B^{lm}_i , \quad (13')$$

where the superscripts u and g mean, respectively, ungerade (odd) and gerade (even) with respect to space inversion. The diagonal component (for simplicity) of the tensor α takes the form

$$\alpha^{(a)}=\frac{1}{12\pi^2}\frac{e^2}{\hbar c}(p_Fl)^2\left[\tau n_i\frac{mp_F}{\pi}\right]^2\frac{2}{3}i\,\mathrm{Im}$$
$$\times\sum_{n_i,i=1,\ldots,4,}\ \sum_{K_i,i=1,\ldots,7}\Lambda^{(a)}(n_i,K_i) , \quad (14)$$

where (a) stands for the representation of the order parameter, and a typical $\Lambda^{(a)}$ is given by

$$\Lambda^{(a)}=A_s^{K_1*}B_a^{K_2K_3}B_a^{K_4K_5*}B_a^{K_6K_7}\int\frac{dO_p\,dO_{p''}\,dO_{p'}}{(4\pi)^3}\hat{\mathbf{p}}\cdot(\hat{\mathbf{p}}'\times\hat{\mathbf{p}}'')$$
$$\times\{\psi_{n_1}^{(u)K_2}(\mathbf{p}')[\psi_{n_2}^{(u)K_4}(\mathbf{p}'')\psi_{n_3}^{(u)K_6}(\mathbf{p})-\psi_{n_2}^{(u)K_4}(\mathbf{p})\psi_{n_3}^{(u)K_6}(\mathbf{p}'')]$$
$$+2\text{ cyclic permutations of }\mathbf{p},\mathbf{p}',\mathbf{p}''\}\ . \quad (15)$$

In this last equation, we have taken for simplicity all $\psi^{(g)}\sim$ identity. Choosing $\psi^{(u)K_2}\sim\hat{\mathbf{p}}_x$, $\psi^{(u)K_4}\sim\hat{\mathbf{p}}_y$, and $\psi^{(u)K_6}\sim\hat{\mathbf{p}}_z$ in the case of the D_{2h} group, we obtain

$$\Lambda^{(a)}\sim\int\frac{dO_p\,dO_{p''}\,dO_{p'}}{(4\pi)^3}[\hat{\mathbf{p}}\cdot(\hat{\mathbf{p}}'\times\hat{\mathbf{p}}'')]^2\ . \quad (16)$$

To conclude this discussion, it is worth mentioning that, while the diagram shown in Fig. 1 gives the main contribution to Eq. (1′), it is proportional to the third power of the order parameter responsible for the violation of space inversion in a new phase.

Before proceeding further, we note that in systems lacking a center of symmetry, an important feature is the so-called photogalvanic effect, i.e., the appearance of a direct current when the sample is uniformly illuminated with light

$$j_i(\omega)=\beta_{iln}(\omega)E_lE_n^*\ . \quad (17)$$

The diagram that provides a finite contribution to the photogalvanic current is shown in Fig. 2. The evaluation of this diagram gives

$$\beta_{\mu\nu\lambda}=i\frac{32\pi^2}{\hbar\omega}\left[\frac{e^2}{\hbar c}\right]^2(p_Fl)^2\frac{\hbar^2c^2}{m}\left[n_i\tau\frac{mp_F}{\pi}\right]^2$$
$$\times\int\frac{dO_p\,dO_{p'}\,dO_{p''}}{(4\pi)^3}\hat{\mathbf{p}}_\mu\hat{\mathbf{p}}'_\nu\hat{\mathbf{p}}''_\lambda W_{\mathbf{pp}'}W^*_{\mathbf{p}'\mathbf{p}}W_{\mathbf{p}'\mathbf{p}''}W^*_{\mathbf{p}''\mathbf{p}'}\ . \quad (18)$$

Because $j_i(\omega)$ is real, the tensor $\beta_{iln}(\omega)$ satisfies $\beta_{iln}=\beta^*_{inl}$. The photogalvanic current can then be rewritten as

$$j_i(\omega)=\beta^s_{iln}E_lE_n^*+i\beta^{as}_{il}(\mathbf{E}\times\mathbf{E}^*)_l\ ,$$

where β^s_{iln} is symmetric with respect to the exchange of the indices l and n. We have

$$\beta^s_{iln}=\tfrac{1}{2}(\beta_{iln}+\beta_{inl})\ ,$$
$$\beta^{as}_{il}=-i\epsilon_{lmn}\beta_{imn}\ , \quad (19)$$

for cubic crystals, $\beta^{as}_{il}=\eta\delta_{ij}$, which leads to the following relation

$$\eta=-i\epsilon_{\mu\nu\lambda}\beta_{\mu\nu\lambda}\ . \quad (20)$$

Therefore, we can write

$$\eta=384\pi^4\frac{e^2}{\hbar c}\frac{\hbar c^2}{m\omega}\alpha\ . \quad (21)$$

The existence of the photogalvanic effect for elliptically polarized light in materials lacking a center of inversion and in the presence of dissipation provides another means to check the results on optical dichroism obtained in certain high-temperature superconductors.

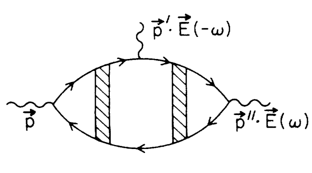

FIG. 2. Diagram giving leading contribution to the photogalvanic coefficient η.

III. OPTICAL DICHROISM ROTATION ANGLE

A. Bulk contribution

We move on now to evaluate explicitly the optical dichroism rotation angle. We consider first the case where the optical dichroism is a bulk effect. Using Eq. (2) and the fact that $\mathbf{B}=\mathbf{H}+4\pi\alpha\mathbf{E}$ for cubic materials we obtain easily the following equation:

$$\nabla\times(\nabla\times\mathbf{E})=i\omega\frac{4\pi}{c^2}\sigma\mathbf{E}+i\omega\frac{4\pi}{c}\alpha\nabla\times\mathbf{E}\ , \quad (22)$$

where σ is the usual conductivity. Writing the incoming wave as $\mathbf{E}_{in}=E_i(\hat{\mathbf{x}}+a\hat{\mathbf{y}})$ and the reflected wave as $\mathbf{E}_{ou}=E_0(\hat{\mathbf{x}}+p_0\hat{\mathbf{y}})$, a right circularly polarized incoming light would correspond to $a=-i$ and a left one to $a=i$. We find that the surface impedance is given by

$$Z=(1-i)\frac{2\pi\omega}{c^2}\left[1+i\alpha\sqrt{\pi\omega/2\sigma}\right] \quad (23)$$

and that

$$p_0\simeq a+\tfrac{1}{2}(1-i)\left\{-2a\sqrt{\pi\omega/2\sigma}+(1-a^2)\frac{2\omega}{\sigma}\right\}\alpha\ . \quad (24)$$

We define the circular dichroism rotation angle between right (+) and left (−) circularly polarized lights by

$$\Phi_{\rm CD}=\frac{|1-ip_0^+|^2+|1-ip_0^-|^2-|1+ip_0^+|^2-|1+ip_0^-|^2}{|1-ip_0^+|^2+|1-ip_0^-|^2}\ . \quad (25)$$

In the limit $\pi\omega\ll\sigma$ (normal skin regime) we find

$$\Phi_{\rm CD}\simeq4\alpha\frac{\omega}{\sigma}\ . \quad (26)$$

For $\omega\sim1\ \mathrm{cm}^{-1}$, $\sigma^{-1}\sim50\ \mu\Omega\,\mathrm{cm}$, we obtain $\Phi_{\rm CD}\simeq14\times10^{-3}$ rad. We note that this result is temperature dependent, since α is temperature dependent. If we assume that $1/\tau\sim T$, we easily find that $\Phi_{\rm CD}\sim1/T$.

B. Anomalous-skin-effect regime

In the case of the anomalous-skin-effect dissipations via impurity or surface roughness scattering are of equal importance. Therefore, we cannot split the two contributions. However, since the relations between ω, $1/\tau$ and δ are not always known, it is sensible to get an estimate for the effect. For simplicity, we consider specular reflection. The equations that we need to solve read

$$\frac{d^2E_x}{dz^2}-i4\pi\alpha\frac{\omega}{c}\frac{dE_y}{dz}=-i\frac{4\pi}{c^2}\omega J_x+2\left[\frac{dE_x}{dz}\right]_{+0}\delta(z)\ ,$$
$$\frac{d^2E_y}{dz^2}+i4\pi\alpha\frac{\omega}{c}\frac{dE_x}{dz}=-i\frac{4\pi}{c^2}\omega J_y+2\left[\frac{dE_y}{dz}\right]_{+0}\delta(z)\ , \qquad (27)$$

where we have chosen the surface of the sample to be the plane $z=0$. Using the solution of this set of equations, we obtain for the quantity p_0 the following value:

$$p_0=a+(1+a)(1+i\sqrt{3})\frac{c^2}{\pi}Z_0^2\alpha\ , \qquad (28)$$

where we have made use of the fact that $Z_0^2\alpha \ll 1$. Here $Z_0=Z_\infty e^{i\pi/3}/2$ and Z_∞ is the surface impedance in the absence of magnetoelectric effect given by[15]

$$Z_\infty=\frac{8}{9}\left[\frac{\sqrt{3}\pi\omega^2 l}{\sigma c^4}\right]^{1/3}(1-i\sqrt{3})\ll 1\ , \qquad (29)$$

where $l=v\tau$ is the mean free path and v is the Fermi velocity. The rotation angle is

$$\Phi_{\rm CD}\simeq\frac{\sqrt{3}}{2}\frac{c^2}{\pi}Z_0^2\alpha\ . \qquad (30)$$

We point out that this result is temperature dependent, since α is temperature dependent. If we assume that $1/\tau\sim T$, we easily find that $\Phi_{\rm CD}\sim 1/T^2$.

C. Infrared regime

We turn now our attention to the case of diffusive boundary conditions and consider the case where $\omega \gg 1/\tau$ and $v/\omega \ll \delta$ (the so-called infrared regime) where δ is the penetration depth. We assume the roughness of the surface to be the source of scattering of the carriers. The current density is given by

$$\mathbf{J}(\omega,z)=-2e\frac{m^3}{h^3}\int\int\int d^3v\,\mathbf{v}f_1\ , \qquad (31)$$

where f_1 is the deviation of the carrier distribution function from the Fermi function f_0. f_1 is a solution to the Boltzmann equation. We impose the following generalized boundary condition

$$f_1^{(2)}(v_x,v_y,v_z,z=0)$$
$$=\int\frac{d\Omega_0}{4\pi}W(\hat{\mathbf{v}},\hat{\mathbf{v}}_0)f_1^{(1)}(v_{0x},v_{0y},-v_{0z},z=0)\ , \qquad (32)$$

where the superscripts 1 and 2 refer, respectively, to before and after the scattering of the carrier with the boundary of the sample. $W(\hat{\mathbf{v}},\hat{\mathbf{v}}_0)$ is the probability of scattering of the carrier on the surface. As in the case of the kinetic magnetoelectric effect, we assume that W is the sum of two parts: a first one, W^s, symmetric, and a second part, W^i, odd with respect to reflections with respect to the plane defined by the projection of the vector $\hat{\mathbf{v}}_0$ on the plane boundary and the z axis. This latter property can be expressed using the azimuthal angles ϕ and ϕ_0 of the cylindrical coordinates as the following

$$W^i(\phi-\phi_0)=-W^i(\phi_0-\phi)\ .$$

For convenience, we are using cylindrical symmetry. In the case of a point group, the reduction in symmetry is that only a single mirror plane perpendicular to the surface is left over. The symmetric part W^s leads to the usual contribution of the diffuse scattering by the surface to the current.[16] In order to estimate the optical dichroism rotation angle we assume the following form for W^i:

$$W^i(\hat{\mathbf{v}},\hat{\mathbf{v}}_0)=\sin(\phi-\phi_0)n_i\ , \qquad (33)$$

where n_i is the fraction of surface defects causing the asymmetric scattering. We obtain the following equations for the components of the electric field

$$\frac{d^2E_x}{dz^2}=i\gamma\int_1^\infty ds\left[\frac{1}{s}-\frac{1}{s^3}\right]\int_0^\infty dz_1E_x(z_1)e^{-i(\omega/v)|z-z_1|s}\ ,$$
$$\frac{d^2E_y}{dz^2}=i\gamma\int_1^\infty ds\frac{1}{s^2}\sqrt{1-1/s^2}e^{-i(\omega/v)zs}\int_1^\infty dt\frac{1}{t}\sqrt{1-1/t^2}\int_0^\infty dz_1E_x(z_1)e^{-i(\omega/v)z_1t}\ , \qquad (34)$$

where

$$\gamma=\frac{1}{4\pi^2}\frac{e^2}{\hbar c}\left[\frac{v}{c}\right]^2\frac{m^2\omega}{\hbar^2}\ .$$

After some calculations we find that the components of the electric field at the surface $z=0$ are related by

$$E_y(0)=\frac{3}{32}n_i\xi E_x(0)\ , \qquad (35)$$

where

$$\xi=\frac{3}{64\pi^3}\left[\frac{v}{c}\right]^2\left[\frac{\omega_p}{\omega}\right]^2\ll 1 \qquad (36)$$

in the infrared regime. In the previous expression ω_p is the plasma frequency. Finally we find that the optical dichroism rotation angle is given by

$$\Phi_{CD} = \frac{9}{128}\frac{n_i}{\pi}\left[\frac{\xi}{2}\right]^{3/4} . \tag{37}$$

We remark that this result is temperature independent. Using $v/c \sim 10^{-3}$ and $m \sim 3m_e$ (these choices are suitable for most high-temperature superconductors) we obtain $\Phi_{CD} \simeq 0.16 n_i / \omega^{(3/2)}$ rad where ω is expressed in units of cm^{-1}.

IV. CONCLUSION

In conclusion, we have presented in this paper an explanation for the optical circular dichroism observed in certain high-temperature superconductors in the normal state. We have shown that for materials lacking a center of inversion as a result of a phase transition (or bulk/surface defects ordering), and in which dissipation plays an important role in phenomena related with reflection or absorbtion of light, the existence of a kinetic magnetoelectric effect provides a natural origin for the occurrence of optical dichroism. We have found a universal relationship between the kinetic magnetoelectric effect (and therefore the circular dichroism rotation angle) and the photogalvanic effect. The simultaneous observation of both optical dichroism and photogalvanic effect will give an important hint concerning the lack of a center of inversion in the materials where the effects are observed either in the bulk or at the surface. It does not, however, give information about the source of this space-inversion symmetry violation.

We also find the optical dichroism rotation angle in both the bulk effect and the anomalous skin effect regimes to be temperature dependent though with a different temperature dependence in each case. In the infrared regime, we find the rotation angle to be temperature independent. This latter behavior is very close to the experimental results in the normal state.[1]

The conflicting data reported for samples of $YBa_2Cu_3O_7$, i.e., the observation or nonobservation of optical dichroism, could be resolved in the following way. It has been shown experimentally that the resistivity of untwinned samples of $YBa_2Cu_3O_7$ is anisotropic[17] in the basal plane. This anisotropy leads to depolarization of reflected light. In some samples, this depolarization could be as big as the optical dichroism rotation angle, and henceforth making the measurement of this latter rather difficult. A second possible explanation to the discrepancy could, of course, be sample dependence of the different experimental results. Finally, one of the motivations behind this work was to show that optical dichroism observed in some high-temperature superconductors can be naturally explained in terms of "conventional" physics without invoking "exotic" physics such as the existence of "anyonic" state of matter.

ACKNOWLEDGMENTS

We would like to acknowledge helpful discussions with D. Pines, D. Ginsberg, M. Salomon, M. Klein, L. R. Testardi, and R. Joynt. This work was supported by National Science Foundation (NSF) Grant No. DMR 89-20538 (B.A.) and the Science and Technology Center for Superconductivity NSF Grant No. DMR 88-09854 (L.G.).

[1]K. B. Lyons, J. Kwo, J. F. Dillon, Jr., G. P. Espinosa, M. McGlashan-Powell, A. P. Ramirez, and L. F. Schneemeyer, Phys. Rev. Lett. **64**, 2949 (1990; K. B. Lyons, J. F. Dillon, Jr., E. S. Hellman, E. H. Hartord, and M. McGlashan-Powell, Phys. Rev. B **43**, 11408 (1991); S. Spielmãn, K. Fesler, C. B. Eom, T. H. Geballe, M. M. Feyer, and A. Kapitulnik, Phys. Rev. Lett. **65**, 123 (1990); H. J. Weber, Th. Wolf, J. Geerk, G. Linker, G. Roth, P. C. Splittgerber-Hunnekes, and G. Guntherodt, Solid State Commun. **76**, 511 (1990).

[2]X. G. Wen and A. Zee, Phys. Rev. B **41**, 240 (1990).

[3]L. D. Landau and E. M. Lifshits, *Electrodynamics of Continuous Media* (Pergamon, New York, 1960).

[4]B. I. Halperin, J. March-Russell, and F. Wilczek, Phys. Rev. B **40**, 8726 (1989); X. G. Wen and A. Zee, Phys. Rev. Lett. **62**, 2873 (1988).

[5]R. J. Kennedy, W. G. Jenks, and L. R. Testardi, Phase Trans. **23**, 19 (1990).

[6]D. Mihailovic and A. J. Heeger, Solid State Commun. **75**, 319 (1990).

[7]D. Mihailovic and I. Poberaj, Physica C **185-189**, 781 (1991).

[8]D. Mihailovic and I. Poberaj (unpublished).

[9]C. Meingast, O. Kraut, T. Wolf, H. Wulh, A. Erb, and G. Muller-Vogt, Phys. Rev. Lett. **67**, 1634 (1991).

[10]I. E. Dzyaloshinskii, Zh. Eksp. Teor. Fiz. **37**, 881 (1959); [Sov. Phys. JETP **10**, 628 (1960)].

[11]W. Opechowski, in *Magnetoelectric Interaction Phenomena in Crystals*, edited by A. J. Freeman and H. Schmid (Gordon and Breach, New York, 1975); I. E. Dzyaloshinskii, Phys. Lett. A **155**, 62 (1991).

[12]L. P. Gor'kov and A. V. Sokol, Pis'ma Zh. Eksp. Teor. Fiz. **45**, 239 (1987) [JETP Lett. **45**, 299 (1987)].

[13]L. S. Levitov, Yu V. Nazarov, and G. M. Eliashberg, Zh. Eksp. Teor. Fiz. **88**, 229 (1985) [Sov. Phys. JETP **61**, 133 (1985)].

[14]V. I. Belinicher and B. I. Sturman, Usp. Fiz. Nauk **130**, 415 (1980) [Sov. Phys. Usp. **23**, 199 (1980)].

[15]C. Kittel, *Quantum Theory of Solids* (Wiley, New York, 1963).

[16]G. E. H. Reuter and E. H. Sondheimer, Proc. R. Soc. London, Ser. A **195**, 336 (1948).

[17]T. A. Friedmann, M. W. Rabin, J. Giapintzakis, J. P. Rice, and D. M. Ginsberg, Phys. Rev. B **42**, 6217 (1990).