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Magneto-optical signature of broken mirror symmetry of two-dimensional conductors

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An unusual aspect of macroscopic electrodynamics of two-dimensional mirror-odd conducting structures bound up with the band spin-orbit coupling $H_{so} = \alpha(\mathbf{p} \times \mathbf{c}) \cdot \boldsymbol{\sigma}$ of current carriers (where $\mathbf{c}$ is one of two nonequivalent normals to a given structure) is pointed out. Namely, it is shown that due to the spin-orbit coupling the presence of the in-plane magnetic field $\mathbf{H}_0$ gives rise to a dependence of the reflection/transmission amplitudes on the structure orientation $\mathbf{c}$, the wave-vector of the incident radiation $\mathbf{q}$, and $\mathbf{H}_0$ of the form $\mathbf{q} \cdot (\mathbf{c} \times \mathbf{H}_0)$. This $\mathbf{q}$- and $\mathbf{H}_0$-odd dependence can be the foundation of the optical way to determine the value of the spin-orbit coupling α.

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I. Introduction. It is well known that substances that not superimpose on their mirror image, for example, when they are comprised of chiral molecules or are non-center-symmetric magnetic crystals, can exhibit unusual physical properties. A remarkable example is the phenomenon of optical activity. In this paper, our aim is to consider a peculiar optical property of another class of materials with destroyed mirror symmetry – two-dimensional (2D) asymmetrical conductors, i.e., 2D conducting structures with broken "up-down" symmetry. The set of such quasi-2D conducting systems having a potential for practical applications in electronics continuously increases. Examples are numerous semiconductor heterostructures with 2D confinement asymmetry [1], as well as the electron gas forming at interfaces of some insulating oxides [2, 3]. A distinguishing property of all conductors mentioned is that the spin-orbit (SO) coupling is an integral part of the dynamics of their current carriers. Namely, since from the symmetry viewpoint the broken "up-down" symmetry of an electron layer is equivalent to the constant electric field piercing the layer, the electron Hamiltonian acquires a term of spin-orbital origin [4]

$$H_{so} = \alpha \, (\mathbf{p} \times \mathbf{c}) \cdot \boldsymbol{\sigma}, \tag{1}$$

where α is a material constant, $\boldsymbol{\sigma}$, $\mathbf{p}$ are, respectively, the Pauli spin and momentum operators, the unit vector $\mathbf{c}$ is one of two nonequivalent normals to the asymmetric structure.

The SO coupling (1) takes part in a great variety of transport phenomena. It not only results in novel mechanisms of some known phenomena such as the spin relaxation [5] and the Anomalous Hall effect [6] but is an initial cause of unusual effects, which would be forbidden in mirror-even systems. Many of them in their nature are anomalous linear or nonlinear responses of a mirror-odd conducting medium to the external electromagnetic field. An example is the magneto-electric effect (MEE) [7, 8], which states that the electric current $\mathbf{J} \sim \mathbf{E}$ passing through the conductor should induce the spin magnetization of the carriers $\mathbf{M} \sim \mathbf{c} \times \mathbf{J}$. The effect inverse to the MEE has also been observed and explained [9]. A derivative of the MEE pointed out recently is that at the presence of the in-plane magnetic field $\mathbf{H}_0$ the current should induce an additional normal component of the spin magnetization $\mathbf{M} \sim \mathbf{c}(\mathbf{H}_0 \cdot \mathbf{J})$ [10]. There is also an experimental indication [11] that the cooperative effect of the dc current $\mathbf{J}$ and the in-plane magnetic field $\mathbf{H}_0$ can be the appearance of the skew-symmetric components of the dynamic conductivity tensor $\delta\sigma_{ij} \sim e_{ijk} c_k (\mathbf{H}_0 \cdot \mathbf{J})$. Another anomalous phenomenon is that the circular polarized electromagnetic wave is capable to induce in the medium the same effects as a constant external magnetic field: (i) the circular polarized light wave (with the electric field $\mathbf{E}_\omega$) induces a permanent spin magnetization $\mathbf{M}_{IF} \sim i\mathbf{c}\,(\mathbf{c} \cdot \mathbf{E}_\omega \times \mathbf{E}^*_\omega)$ [13], which means the inverse Faraday effect, and (ii) if the constant electric current $\mathbf{J}$ passes through the system, the circular wave induces the Hall-like current $\mathbf{J}_{\mathrm{H}} \sim i(\mathbf{c} \times \mathbf{J})(\mathbf{c} \cdot \mathbf{E}_\omega \times \mathbf{E}^*_\omega) \sim \mathbf{M}_{IF} \times \mathbf{J}$ [14]. Mutual influence of charge and spin flows due to H_{so} was also considered (see, e.g., Ref. [8] and references therein).

But besides that, there are direct consequences of destroyed mirror symmetry of quite a different kind.

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These are phenomena associated with the invariant $\mathbf{q}\times\times(\mathbf{c}\times\mathbf{H}_0)$ which is forbidden in P-even systems. The feature of this invariant is that it is the odd function of the applied magnetic field $\mathbf{H}_0$, the orientation of the system $\mathbf{c}$, and the vector $\mathbf{q}$, which in situations studied earlier was the wave-vector of some excitations in a mirror-odd media. For the first time the invariant has appeared in the physics of excitons in crystals of polar symmetry to explain the linear in $\mathbf{H}_0$ energy shift of an exciton in CdS moving with the momentum $\mathbf{q}$ – the Magneto-Stark effect [15]. (The symmetry, the 2D structures considered in this paper is a 2D analogue of 3D polar symmetry.) Later on, it was experimentally shown that the invariant is responsible for the $\mathbf{q}$-odd dependence of the damping of the spin-density-wave with the wave-vector $\mathbf{q}$ in doped CdS [16]. In polar-symmetry superconductors subject to the weak magnetic field $\mathbf{H}_0$, the energy of the Cooper pair with the momentum $\mathbf{q}$ was shown to include a term proportional to $\mathbf{q}\cdot(\mathbf{c}\times\mathbf{H}_0)$ which results in a phase of the Cooper-pair condensate [17]. Three examples mentioned of the unusual behavior of Bose-type excitations make it natural to expect that the electromagnetic field may also acquire an unusual property due to the interaction with a mirror-odd medium. In particular, one may expect that some characteristics of an electromagnetic wave could depend on the invariant $\mathbf{q}\cdot(\mathbf{c}\times\mathbf{H}_0)$, where $\mathbf{q}$ is the wave-vector of the wave. Since the reflecting wave is specified by the reflection and transmission amplitudes, this is equivalent to a dependence of these amplitudes on the invariant, where $\mathbf{q}$ is the wave-vector of the incident radiation. It is shown below that this proposal is true.

The subject under study has also a practical aspect. Although the SO coupling (1) plays an essential role in many recent achievements in the physics of spin-dependent phenomena [6, 18] and the value of α is very important for their practical employments, there are still few experimental methods that can be used to determine it in a structure given. The first method was an analysis of the beat patterns observed in Shubnikov–de Haas oscillations at hight magnetic fields [19]. This method requires a very high purity of samples in order to observe tens of the oscillations. There is also the method of angle-resolved photoemission spectroscopy, which was used to determine α from the dispersion of some surface states [20]. This paper suggests an optical method – the measurement of the magnetic-field dependence of the reflection/transmission amplitudes.

II. Constitutive relations. Within the frame of linear electrodynamics, the magnetic field can have an effect on the wave only through the so-called material constitutive relations [21], which express the electric current $\mathbf{J}$ and the electron magnetization $\mathbf{M}$ through the electric and magnetic fields, $\mathbf{E}$ and $\mathbf{H}$. The general form of these relations in asymmetric 2D conductors is given by [22]

$$\mathbf{M}=\hat{\chi}\mathbf{B}+\hat{\gamma}\mathbf{E}, \tag{2}$$

$$\mathbf{J}=\hat{\sigma}\mathbf{E}-\hat{\theta}\dot{\mathbf{H}}, \tag{3}$$

where a point over $\mathbf{H}$ denotes the time derivative. The presence of the cross terms is a specific property of a conducting media of polar symmetry; two terms in the right hand site of the equations have opposite parity under space inversion. The magneto-electric susceptibilities $\hat{\gamma}$ and $\hat{\theta}$ are material characteristics additional to the electric conductivity $\hat{\sigma}$ and the magnetic susceptibility $\hat{\chi}$. Their explicit form was earlier calculated at the absence of magnetic field and at the magnetic field perpendicular to the 2D conductor at frequencies near the spin-resonance frequency [22]. Here we consider the in-plane magnetic field $\mathbf{H}_0$ and show that at sufficiently high frequencies of the incident radiation ω the main contribution to the tensors $\hat{\gamma}$ and $\hat{\theta}$ are linear in $\mathbf{H}_0$ and have the form $\gamma_{ij}\sim\alpha c_i H_{0\,j}$ and $\theta_{ij}\sim\alpha H_{0\,i}c_j$ and then derive the Fresnel formulas showing the sought-for effect.

The Hamiltonian of the system under consideration is

$$H=H_0+H_{so}+H_{\mathrm{Z}}, \tag{4}$$

where $H_0=\mathbf{p}^2/2m+H_{\mathrm{imp}}$, $\sum_i U\delta(\mathbf{r}-R_i)$ is the interaction with arbitrary distributed short-range impurities, $H_{\mathrm{Z}}=-\mu\cdot\mathbf{H}_0$ is the Zeeman interaction, and H_{so} is given by (1). There are several dimensionless parameters determining the kinetics of the system. It is assumed that the parameter $\epsilon_{\mathrm{F}}\tau$, where $\epsilon_{\mathrm{F}}=\mathbf{p}_{\mathrm{F}}^2/2m$ is the Fermi energy and τ is the mean scattering time given by $\tau^{-1}=mn_{\mathrm{imp}}|U|^2$, is large. The applied magnetic field appears in the theory through the parameter $\Omega_s=\omega_s\tau$, where $\omega_s=|g|\mu_{\mathrm{B}}H_0$. It is treated as small. The SO constant α appears in the problem in two ways. The parameter $\delta=\alpha p_{\mathrm{F}}/\epsilon_{\mathrm{F}}$ is treated as very small and will be ignored. Another parameter $\eta=2\alpha p_{\mathrm{F}}\tau$ controls the kinetics of spin-flip processes by impurity scattering. It is also treated as small but all necessary powers of η will be taken into account. The frequency of the electromagnetic field ω enters through the parameter $\omega\tau$ which is considered as moderately large $1<\omega\tau<\epsilon_{\mathrm{F}}\tau$. Consider $\hat{\gamma}$ first. Its microscopic definition is given by [22] $\gamma_{ij}(\omega)=\frac{e}{\omega}\left(\frac{g\mu_{\mathrm{B}}}{2}\right)\int_0^\infty dte^{i(\omega+i0)t}\langle[\sigma_i(t),v_j]\rangle$. The evaluation of this retarded correlation function at $\epsilon_{\mathrm{F}}\tau\gg 1$ is known to reduce to the summation of the ladder diagrams shown in Fig. 1. By applying standard methods

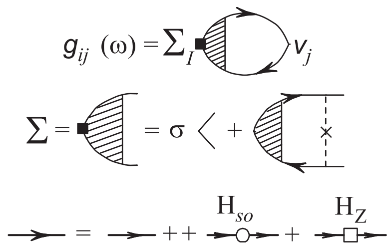

Fig. 1. Diagrammatic representation for the spin-velocity correlation function $\gamma_{ij}(\omega)$ and Bethe–Salpeter equation satisfied by the renormalized spin vertex $\mathbf{\Sigma}_{i,\,\alpha\beta}(i\omega)$. The thick solid line represents the single-particle Green's function $G_{\alpha,\beta}(\mathbf{p}|\,i\epsilon)$, the thin solid line represents the bar Green's function corresponding to H_0, and the dashed line with a cross on it corresponds to the correlator of random potential

of statistical mechanics [23, 24], we have at zero temperature

$$\gamma_{ij}(\omega) = m\tau\left(\frac{eg\mu_{\rm B}}{4\pi}\right)\mathrm{Tr}\left\{\Sigma^i v^j_{(1)}(\omega)\right\}. \qquad (5)$$

Here $\mathbf{\Sigma}_{\alpha\beta}(\omega)$ is the impurity-renormalized spin vertex and $\mathbf{v}^{(1)}_{\kappa\beta}(\omega) = \frac{1}{m\tau}\int_{\mathbf{p}}\left[G^R(\epsilon_{\rm F},\mathbf{p})\mathbf{v}(\mathbf{p})G^A(\epsilon_{\rm F}+\omega,\mathbf{p})\right]_{\kappa\beta}$, where $G^{R(A)}$ is the retarded (advanced) Green function corresponding to the Hamiltonian (4) and $\int_{\mathbf{p}} = \int d^2p/(2\pi)^2$, has the sense of the first impurity correction to the bare velocity operator. The evaluation of $\hat{\mathbf{v}}^{(1)}(\omega)$, $\hat{\Sigma}(\omega)$, and $\gamma_{ij}(\omega)$ can be performed following the method of Refs. [13, 20] which allows to obtain solutions to diagrammatic equations for any relevant quantities without making uncontrolled assumptions about their spin-matrix structure. The point-like impurity potential allows one to represent the Bethe–Salpeter equation for the spin vertex $\mathbf{\Sigma}$ as a matrix equation

$$\mathbf{\Sigma}_{\alpha\beta}(\omega) = \boldsymbol{\sigma}_{\alpha\beta} + \mathbf{\Sigma}_{\kappa\delta}(\omega)T_{\delta\kappa|\alpha\beta}(\omega), \qquad (6)$$

where the kernel is given by $T(\omega)_{\delta\kappa|\alpha\beta} = \frac{1}{m\tau}\int_{\mathbf{p}} G^R_{\delta\beta}(\epsilon_{\rm F}+\omega,\mathbf{p})G^A_{\alpha\kappa}(\epsilon_{\rm F},\mathbf{p})$. By expanding the electron Green function into series in H_{so} and $H_{\rm Z}$, as it is shown in Fig. 1, one gets the explicit form of $\hat{\mathbf{v}}^{(1)}$ and $\hat{T}(\omega)$ as power series about the parameters η and Ω_s

$$\hat{\mathbf{v}}^{(1)}(\omega) = -\alpha\left[\frac{1}{2}\eta_\omega f_\omega(\mathbf{c}\times\boldsymbol{\sigma}) + \Omega_s f_\omega^2\,\mathbf{h}(\mathbf{c}\cdot\boldsymbol{\sigma})\right], \qquad (7)$$

where $f_\omega = (1-i\Omega)^{-1}$, $\Omega = \omega\tau$, $\mathbf{h} = \mathbf{H_0}/|\mathbf{H_0}|$, $\eta_\omega = \eta f_\omega$, and

$$T(\omega) = T^{(0)}(\omega) + T^{(1)}_{(\rm Z)}(\omega) + T^{(2)}_{(so)}(\omega), \qquad (8)$$

where

$$T^{(0)}_{\delta\kappa|\alpha\beta}(\omega) = \delta_{\delta\beta}\delta_{\alpha\kappa}f_\omega,$$

$$T^{(1)}_{(Z)\,\delta\kappa|\alpha\beta} = \frac{i}{2}\Omega_s f_\omega^2\left[(\mathbf{h}\cdot\boldsymbol{\sigma})_{\delta\beta}\delta_{\alpha\kappa} - \delta_{\delta\beta}(\mathbf{h}\cdot\boldsymbol{\sigma})_{\alpha\kappa}\right], \qquad (9)$$

$$T^{(2)}_{(so)\,\delta\kappa|\alpha\beta} = \frac{\eta^2}{4}f_\omega^3\left[(\mathbf{c}\times\boldsymbol{\sigma})_{\delta\beta}\cdot(\mathbf{c}\times\boldsymbol{\sigma})_{\alpha\kappa} - 2\delta_{\delta\beta}\delta_{\alpha\kappa}\right].$$

Following Refs. [16, 23], one ought to solve Eq. (6) with the kernel (8) by making use of matrixes $\delta_{\alpha\beta}$, $(\mathbf{c}\cdot\boldsymbol{\sigma})_{\alpha\beta}$, $(\mathbf{h}\cdot\mathbf{c}\times\boldsymbol{\sigma})_{\alpha\beta}$, and $[(\mathbf{c}\times\mathbf{h})\cdot(\mathbf{c}\times\boldsymbol{\sigma})]_{\alpha\beta}$ as a basis in the spin-matrix space. The result is

$$\mathbf{\Sigma}_{\alpha\beta}(\omega) = \mathbf{P}(\mathbf{c}\cdot\boldsymbol{\sigma})_{\alpha\beta} + \mathbf{Q}(\mathbf{h}\cdot\mathbf{c}\times\boldsymbol{\sigma})_{\alpha\beta} + \mathbf{R}[(\mathbf{c}\times\mathbf{h})\cdot(\mathbf{c}\times\boldsymbol{\sigma})]_{\alpha\beta}, \qquad (10)$$

where

$$\mathbf{P} = \mathbf{c}(-i\Omega + \eta_\omega^2/2)(f_\omega Z)^{-1} - (\mathbf{c}\times\mathbf{h})\Omega_s Z^{-1},$$
$$\mathbf{Q} = \mathbf{c}\Omega_s Z^{-1} + (\mathbf{c}\times\mathbf{h})(-i\Omega + \eta_\omega^2)(f_\omega Z)^{-1}, \qquad (11)$$
$$\mathbf{R} = \mathbf{h}(-i\Omega + \eta_\omega^2/2)^{-1}f_\omega^{-1},$$
$$Z = -\Omega^2 - i\cdot\frac{3}{2}\Omega\eta_\omega^2 + \frac{1}{2}\eta_\omega^4 + \Omega_s^2 f_\omega^2.$$

The substitution of (7) and (10) into Eq. (5) yields

$$\gamma_{ij} = \alpha m\tau\,(eg\mu_{\rm B})\,f_{ij}, \qquad (12)$$
$$f_{ij} = Ac_ih_j + B(\mathbf{c}\times\mathbf{h})_i h_j - Ch_i(\mathbf{c}\times\mathbf{h})_j, \qquad (13)$$

where $A = i\Omega\Omega_s f_\omega Z^{-1}$, $B = \frac{\eta_\omega^2}{2}(-i\Omega + \eta_\omega^2)Z^{-1}$, $C = \frac{\eta_\omega^2}{2}\left(-i\Omega + \frac{\eta_\omega^2}{2}\right)^{-1}$. Although Eqs. (9) are derived at $\omega_s\tau < 1$, they remain qualitatively valid also at $\omega_s\tau \sim 1$ up to corrections of relative order of $(\omega_s\tau)^2$. The corrections do not change the matrix structure of (10) and hence are of no importance[2)]. An analogous consideration of $\hat{\theta}$, whose microscopic definition is $\theta_{ij}(\omega) = \frac{e}{\omega}\left(\frac{g\mu_{\rm B}}{2}\right)\int_0^\infty dt e^{i(\omega+i0)t}\langle[v_i(t),\sigma_j]\rangle$, shows that $\theta_{ij} = \gamma_{ij}(\mathbf{h}\to-\mathbf{h})$. At $\eta < 1$ and $\omega\tau > 1$ so that $\eta^2/(\omega\tau)^2 \ll 1$, the first term in (13) dominates. Thus the constitutive relations (2), (3) take the form

$$\mathbf{J}_\omega = \sigma_\omega\mathbf{E}_\omega - i\omega\underline{\kappa}\mathbf{H_0}(\mathbf{c}\cdot\mathbf{H}_\omega), \qquad (14)$$
$$\mathbf{M}_\omega = \underline{\kappa}\mathbf{c}(\mathbf{H_0}\cdot\mathbf{E}_\omega), \qquad (15)$$

where $\underline{\kappa} = \kappa_\omega|\mathbf{H_0}|^{-1}$, $\kappa_\omega = \alpha m\tau(eg\mu_{\rm B})A$ and the first term in Eq. (2) was dropped in view of very small value of the spin susceptibility χ. In the following the current and magnetization obtained will be provided

[2)] In the limit $\omega\tau\to 0$ the magnetic field dependence disappears from γ_{ij} in accordance with Ref. [10] and, on account of the equality $h_i(\mathbf{c}\times\mathbf{h})_j - (\mathbf{c}\times\mathbf{h})_i h_j = e_{ijn}c_n$, we get the known result [7] $\mathbf{M}\sim\mathbf{c}\times\mathbf{E}$.

with the subscript S showing their 2D character so that corresponding 3D densities are $\mathbf{J} = \delta(z)\mathbf{J}_S$ and $\mathbf{M} = \delta(z)\mathbf{M}_S$.

III. Fresnel formulas. Let the quasi-2D layer aligned along an $x{-}y$ plane is placed at position $z = 0$ between two dielectrics with the permittivity ϵ_1 $(z < 0)$ and ϵ_2 $(z > 0)$, and z-axis points "upward" to the dielectric 2. By excluding the magnetic field $\mathbf{H}_\omega$ from the Maxwell equations

$$\nabla \times \mathbf{E}_\omega = iq_0 \left(\mathbf{H}_\omega + 4\pi\mathbf{M}_\omega\right), \tag{16}$$

$$\nabla \times \mathbf{H}_\omega = -iq_0\epsilon\mathbf{E}_\omega - \frac{4\pi}{c}\mathbf{J}_\omega, \tag{17}$$

where $q_0 = \omega/c$, one can obtain the conventional wave equation for the electric field $\mathbf{E}_\omega$. The electric field can be represented as

$$\mathbf{E}_\omega(\mathbf{r}) = e^{i\mathbf{q}\cdot\mathbf{r}}\left[\mathbf{E}(z),\, E_3(z)\right], \tag{18}$$

where $\mathbf{q} = (q_1,\, q_2,\, 0)$ and $\mathbf{E}(z) = (E_1(z),\, E_2(z),\, 0)$ are parallel to the plane. Then the wave equation outside the layer falls into two equations

$$\left[\partial_z^2 + p^2(z)\right]\mathbf{E}(z) = 0, \tag{19}$$

$$i\partial_z\left[\mathbf{q}\cdot\mathbf{E}(z)\right] - p^2(z)E_3(z) = 0, \tag{20}$$

where $p^2(z) = \epsilon(z)q_0^2 - \mathbf{q}^2$ and $\epsilon(z) = \theta(-z)\epsilon_1 + \theta(z)\epsilon_2$. From the Faraday law it follows

$$\hat{\mathbf{n}} \times (\mathbf{E}_{\omega,\,2} - \mathbf{E}_{\omega,\,1}) = 4\pi iq_0\mathbf{M}_{\omega,\,S\|}, \tag{21}$$

where $\mathbf{n} = \hat{\mathbf{z}}$, $\mathbf{E}_{\omega,\,1}$ and $\mathbf{E}_{\omega,\,2}$ are the values of the electric field on the lower and upper sides of the layer, respectively, and $\mathbf{M}_{\omega,\,S\|}$ is the parallel component of the 2D magnetization density. According to Eq. (15), $\mathbf{M}_{\omega,\,S\|} = 0$ consequently $\mathbf{E}(z)$ has not a jump on the layer. The boundary conditions which follow from the Amper law have the conventional form

$$\hat{\mathbf{n}} \times (\mathbf{H}_{\omega,\,2} - \mathbf{H}_{\omega,\,1}) = \frac{4\pi}{c}\mathbf{J}_{\omega,\,S}, \tag{22}$$

where $\mathbf{J}_{\omega,\,S}$ is given by Eq. (14). Note that Eq. (22) specifying the jump of parallel components of the magnetic field leaves its perpendicular component continuous. Therefore the expression (14) for the electric current remains well defined. Using the Maxwell equations and the constitution relations, Eq. (22) can be put in the form

$$\Delta_{ij}^{(2)}\partial_z E_j(0^+) - \Delta_{ij}^{(1)}\partial_z E_j(0^-) = -iD_{ij}E_j(0), \tag{23}$$

where $i, j = x, y$, $\Delta_{ij}^{(1,\,2)} = \left(\delta_{ij} + q_iq_j/p_{(1,\,2)}^2\right)$, $p_{(1,\,2)}^2 = q_0^2\epsilon_{1,\,2} - \mathbf{q}^2$, and $D_{ij} = \frac{4\pi}{c}q_0\left[\sigma\delta_{ij} - i\kappa c h_i(\mathbf{c}\times\mathbf{q})_j\right]$.

For the wave incident from the dielectric 1 the solution of Eq. (19) can be sought in the form

$$\mathbf{E}(z) = \theta(-z)\left(\mathbf{a}e^{ip_1z} + \mathbf{b}e^{-ip_1z}\right) + \theta(z)\mathbf{t}e^{ip_2z}, \tag{24}$$

where $\mathbf{a}$, $\mathbf{b}$, and $\mathbf{t}$ are 2D vectors. The perpendicular component $E_3(z)$ is obtained from Eq. (19) so that the total electric fields of incident $\mathbf{E}_{(in)}$, reflected $\mathbf{E}_{(re)}$, and transmitted $\mathbf{E}_{(tr)}$ waves are

$$\begin{gathered}\mathbf{E}_{(in)}(\mathbf{r}) = e^{i\mathbf{q}_{(in)}\cdot\mathbf{r}}\left[\mathbf{a}, \frac{-1}{p_1}(\mathbf{q}\cdot\mathbf{a})\right],\\ \mathbf{E}_{(re)}(\mathbf{r}) = e^{i\mathbf{q}_{(re)}\cdot\mathbf{r}}\left[\mathbf{b}, \frac{1}{p_1}(\mathbf{q}\cdot\mathbf{b})\right],\\ \mathbf{E}_{(t)}(\mathbf{r}) = e^{i\mathbf{q}_{(t)}\cdot\mathbf{r}}\left[\mathbf{t}, \frac{-1}{p_2}(\mathbf{q}\cdot\mathbf{t})\right],\end{gathered} \tag{25}$$

where $\mathbf{q}_{(in)} = (\mathbf{q}, \mathbf{n}p_1)$, $\mathbf{q}_{(re)} = (\mathbf{q}, -\mathbf{n}p_1)$, and $\mathbf{q}_{(tr)} = (\mathbf{q}, \mathbf{n}p_2)$ are the wave vectors of the incident, reflected, and transmitted waves, respectively. The continuity at $z = 0$ and the boundary conditions (23) yield

$$\begin{gathered}t_i = 2p_1\left[(p_1 + p_2)\cdot\hat{1} + i\hat{F}\right]_{ij}^{-1} a_j,\\ r_i = \left\{\left[(p_1 - p_2)\cdot\hat{1} - i\hat{F}\right]\left[(p_1 + p_2)\cdot\hat{1} + i\hat{F}\right]^{-1}\right\}_{ij} a_j,\\ F_{ij} = Ne_i^{(1)}e_j^{(1)} + Me_i^{(2)}e_j^{(2)} + Le_i^{(1)}e_j^{(2)},\end{gathered} \tag{26}$$

where $\mathbf{e}^{(1)} = \mathbf{q}/|\,\mathbf{q}|$ and $\mathbf{e}^{(2)} = (\mathbf{c}\times\mathbf{q})/|\,\mathbf{q}|$ form the basis in x, y-plane bound to the wave-vector of the incident radiation,

$$\begin{gathered}N = -iuq_0\left(1 - \frac{\mathbf{q}^2}{\epsilon_1q_0^2}\right) - i\frac{\epsilon_1 - \epsilon_2}{\epsilon_1p_2}\mathbf{q}^2,\\ M = -v(\mathbf{h}\cdot\mathbf{c}\times\mathbf{q}) - iuq_0,\\ L = -v(\mathbf{h}\cdot\mathbf{q})\frac{p_1^2}{\epsilon_1q_0^2},\end{gathered} \tag{27}$$

and $u = 4\pi\sigma_\omega/c$, $v = 4\pi\kappa_\omega q_0$.

Eqs. (26) show that the transmitted and reflected waves are depolarized in the general case. The effect of the invariant $\mathbf{q}\times\mathbf{c}\cdot\mathbf{H}_{(0)}$ is more pronounced for the s-polarized incident radiation, i.e., at $\mathbf{a}_{(s)} = \mathbf{n}\times\mathbf{q}$. In this case, if one directs $\mathbf{H}_{(0)}$ perpendicular to the plane of incidence, the reflected and transmitted waves are also s-polarized and the Fresnel formulas have the form

$$\mathbf{r}_{(s)} = \frac{n_1\cos\theta_1 - n_2\cos\theta_2 - \frac{4\pi\sigma_\omega}{c}\left(1 - i\frac{\omega\kappa_\omega}{\sigma_\omega}\frac{\mathbf{h}\cdot\mathbf{c}\times\mathbf{q}}{q_0}\right)}{n_1\cos\theta_1 + n_2\cos\theta_2 + \frac{4\pi\sigma_\omega}{c}\left(1 - i\frac{\omega\kappa_\omega}{\sigma_\omega}\frac{\mathbf{h}\cdot\mathbf{c}\times\mathbf{q}}{q_0}\right)}\mathbf{a}_{(s)}, \tag{28}$$

$$\mathbf{t}_{(s)} = \frac{2n_1\cos\theta_1}{n_1\cos\theta_1 + n_2\cos\theta_2 + \frac{4\pi\sigma_\omega}{c}\left(1 - i\frac{\omega\kappa_\omega}{\sigma_\omega}\frac{\mathbf{h}\cdot\mathbf{c}\times\mathbf{q}}{q_0}\right)}\mathbf{a}_{(s)}, \quad (29)$$

where $n_{1,2} = \sqrt{\epsilon_{1,2}}$ and $\theta_{1,2}$ are, respectively, the angles of incidence and refraction.

IV. Summary. A characteristic feature of magneto-optics of 2D mirror-odd conductors is found, namely if the conductor is subject to the in-plane magnetic field $\mathbf{H}_0$, the reflection/transmission amplitudes contain $\mathbf{H}_0$-odd terms proportional to the P- and T-odd invariant $\mathbf{c}\cdot(\mathbf{q}\times\mathbf{H}_0)$, where $\mathbf{q}$ is the wave-vector of the incident wave and $\mathbf{c}$ is one of two nonequivalent normals to the conductor. It should be stressed that this dependence characterizes the very electromagnetic *field*. A *material* property with a similar optical manifestation might be the presence of terms bilinear in vectors $\mathbf{c}\times\mathbf{q}$, and $\mathbf{H}_0$ in the conductivity tensor σ_{ij} on account of the space dispersion. However, a straightforward but lengthy calculation shows that within the scope of the model considered (the quadratic energy spectrum of electrons and the s-wave impurity scattering) such terms cancel each other in the α-linear approximation[3)]. The effect found demonstrates the violation of the reciprocity theorem in the system – the reflection of the wave with $\mathbf{q}$ differs from that with $-\mathbf{q}$. In addition, it gives a way to determine the value of this coupling in 2D asymmetric structures. The factor $\frac{\omega\kappa_\omega}{\sigma_\omega} \sim g\frac{\alpha}{v_F}\frac{m}{m_0}\frac{v_F}{c}(\omega_s\tau)$, where m_0 is the electron mass in vacuum, is small – it is of the order of 10^{-5} for an ordinary metallic monolayer and can be more for a monolayer composed of heavy atoms, e.g., transition metal oxides [2] and heavy fermion compounds [26]. Therefore, only terms linear in α should be kept. However, the background signal can be suppressed in favor of the SO-induced corrections if one takes advantage of the fact that the α-linear corrections are simultaneously linear in $\mathbf{H}_0$. This circumstance allows to apply the modulation method, when one modulates $\mathbf{H}_0$ in amplitude with some low frequency and detects the variation of reflected/transmitted wave on the modulation frequency. The described way of determination of α has the following advantages. (i) It does not suggest electrical measurements, hence, there is no need for contacts. (ii) The conducting layer under study may not have a contact with vacuum, being inside the sample. (iii) There are not heavy restrictions on the idealness of the conducting layer. The effect described is not peculiar to 2D systems; the reflection off metals of polar symmetry should also possess an analogous feature.

[3)] In addition, the wave vector $\mathbf{q}$ enters those bilinear terms through the dimensionless combination $\mathbf{q}v_F\tau$, which is less than the combination $\mathbf{q}/q_0$ obtained from the Maxwell equations.

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