Chiral Magnetic Photocurrent in Dirac and Weyl Materials

Sahal Kaushik,[1,*] Evan John Philip,[1,†] and Dmitri E. Kharzeev[1,2,3,‡]

[1] *Department of Physics and Astronomy, Stony Brook University, Stony Brook, New York 11794-3800, USA*
[2] *Department of Physics, Brookhaven National Laboratory, Upton, New York 11973-5000, USA*
[3] *RIKEN-BNL Research Center, Brookhaven National Laboratory, Upton, New York 11973-5000, USA*

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Circularly polarized light (CPL) can induce the asymmetry between the number of left- and right-handed chiral quasiparticles in Dirac and Weyl semimetals. We show that if the photoresponse of the material is dominated by chiral quasiparticles, the total chiral charge induced in the material by CPL can be evaluated in a model-independent way through the chiral anomaly. In the presence of an external magnetic field perpendicular to the incident CPL, this allows to predict the linear density of the induced photocurrent resulting from the chiral magnetic effect. The predicted effect should exist in any kind of Dirac or Weyl semimetals, with both symmetric and asymmetric band structure. An estimate of the resulting *chiral magnetic photocurrent* in a typical Dirac semimetal irradiated by an infrared laser of intensity $\sim 5 \times 10^6\,\mathrm{W/m^2}$ and a wavelength of $\lambda \simeq 10\,\mu\mathrm{m}$ in an external magnetic field $B \sim 1\,\mathrm{T}$ yields an extremely large current $J \sim 20\,\mu\mathrm{A}$ in the laser spot of size $\sim 50\,\mu\mathrm{m}$. This opens possibilities for applications in photonics, optoelectronics, and THz sensing.

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Circularly polarized light (CPL) breaks the symmetry between left and right and thus possesses a non-zero chirality. In the interactions of CPL with a material, the chirality of light can be transferred to the chirality of matter. Of particular interest is the interaction of CPL with the 3-dimensional chiral quasiparticles in recently discovered Dirac and Weyl semimetals[1–5]. In this note, we will show that the chirality transfer from light to fermions can be described in a model-independent way by using the chiral anomaly[6,7].

The physical effect of the chiral anomaly is the following. If the charged chiral fermion quasiparticles are massless (the corresponding linear band has no gap, in the context of condensed matter physics), and no chirality-changing interactions are present in the Hamiltonian, the chirality of external gauge fields can be transferred to the chirality carried by the charged quasiparticles, and viceversa. The only conserved quantity is the total chirality of the fermion quasiparticles and the gauge field.

Apart from the transfer of chirality from CPL, the chiral anomaly facilitates the transport of charge in parallel electric and magnetic fields by creating the chirality imbalance[8,9]. The resulting longitudinal negative magnetoresistance[10,11] has been observed in Dirac semi-metals such as $ZrTe_5$[12] and Na_3Bi[13] and Weyl semi-metals such as TaAs[14]. Recently, a large photocurrent due to CPL has been observed in TaAs, a Weyl semimetal with tilted cones, in which the cancelation of the photocurrent between the Weyl cones can be avoided[15]. The effect that we will discuss does not rely on the tilt of the cones and is distinct. CPL has also been proposed to cause a Photovoltaic Hall effect[16,17] in graphene, which is a 2-dimensional material with relativistic Dirac fermion quasiparticles but, as we will see, the underlying mechanism is different in 3D materials due to the chiral anomaly. The chiral pumping in 3D Dirac materials by rotating electric fields has been considered in[18].

In the interaction of the CPL with an optically thick Dirac or Weyl semimetal (for a mid-infrared laser, the light penetration length for these materials is of the order of a few hundred nanometers), the chirality of the absorbed light gets fully transferred to the material. If the optical response of the material is dominated by the chiral quasi-particles (e.g. when other electronic bands cannot be excited at a given light frequency), the chirality of the CPL gets tranferred to the asymmetry between the number of left- and right-handed chiral quasiparticles. The resulting chiral asymmetry, as we will see, is completely fixed by the chiral anomaly. The chiral asymmetry in an external magnetic field is known to induce the chiral magnetic effect, the magnitude of which is also fixed by the chiral anomaly. Therefore, the use of chiral anomaly allows us to evaluate the magnitude of the resulting chiral magnetic photocurrent.

The chirality carried by the chiral fermion quasiparticles is described by the axial current $j_5^\mu = \bar{\psi}\gamma^\mu\gamma^5\psi$, where ψ is the quasiparticle's spinor wave function; the temporal component of this current is the density of chiral charge, $j_5^0 \equiv \rho_5$.

The chiral anomaly causes non-conservation of j_5^μ in the presence of an electromagnetic field $F^{\mu\nu}$, as given by[6,7]

$$\partial_\mu j_5^\mu = \frac{e^2}{16\pi^2}\epsilon^{\mu\nu\rho\sigma}F_{\mu\nu}F_{\rho\sigma}; \quad (1)$$

note that the quantity on the right hand side is odd under parity and thus vanishes for linearly polarized light. It is given by the full derivative of the Chern-Simons current

$$h^\mu = -\frac{e^2}{8\pi^2}\epsilon^{\mu\nu\rho\sigma}A_\nu F_{\rho\sigma} \quad (2)$$

that describes the chirality density and flux carried by the electromagnetic field. Therefore, as follows from (1) and (2), the sum of the chirality of the electromagnetic field and the chirality of the fermions is conserved:

$$\partial_\mu(j_5^\mu + h^\mu) = 0. \quad (3)$$

The chirality density and chirality flux are given (in SI units) by

$$h^0 = \frac{e^2}{4\pi^2\hbar^2}\boldsymbol{A}\cdot\boldsymbol{B} \tag{4}$$

and

$$\boldsymbol{h} = \frac{e^2}{4\pi^2\hbar^2}(A^0\,\boldsymbol{B} - \boldsymbol{A}\times\boldsymbol{E}). \tag{5}$$

The chirality density h^0 is known in magneto-hydrodynamics, where *magnetic helicity*[19–22] is defined as $\int \boldsymbol{A}\cdot\boldsymbol{B}\,\mathrm{d}^3r$. In our case, the prefactor of $e^2/(4\pi^2\hbar^2)$ appears in equation (4) because, as it can be seen from equation (3), it is the chirality density available for the transfer to the chiral fermions.

In a real Dirac or Weyl material, there is also chirality-flipping scattering, with a characteristic relaxation time τ_{V}, so equation (1) (in SI units) becomes

$$\dot{\rho_5} + \nabla\cdot\boldsymbol{j}_5 = \frac{e^2}{2\pi^2\hbar^2}\boldsymbol{E}\cdot\boldsymbol{B} - \frac{\rho_5}{\tau_{\mathrm{V}}}. \tag{6}$$

The last term on the right hand side of equation (6) does not affect the balance of chirality transfer if the frequency of light ω is large compared to ${\tau_{\mathrm{V}}}^{-1}$.

For an oscillating electric field with $\boldsymbol{E}(t,\boldsymbol{r}) = \Re(\mathrm{e}^{-\mathrm{i}\omega t}\boldsymbol{\mathcal{E}}(\boldsymbol{r}))$, in the Coulomb gauge with $A^0 = 0$, the time-averaged chirality flux (5) is

$$\langle\boldsymbol{h}\rangle = \frac{e^2}{8\pi^2\hbar^2\omega}\Re(\mathrm{i}\,\boldsymbol{\mathcal{E}}\times\boldsymbol{\mathcal{E}}^*). \tag{7}$$

For light traveling in the z direction, this becomes

$$\langle h^z\rangle = \frac{e^2}{4\pi^2\hbar^2\omega}\Re(\mathrm{i}\,\mathcal{E}_x\mathcal{E}_y^*); \tag{8}$$

note that this quantity vanishes for linearly polarized light.

For CPL with $\boldsymbol{\mathcal{E}} = E_0\,(\hat{\boldsymbol{x}} \pm \mathrm{i}\,\hat{\boldsymbol{y}})$,

$$\langle h^z\rangle = \pm\frac{e^2}{4\pi^2\hbar^2\omega}E_0^2. \tag{9}$$

The ratio of this chirality flux to the energy flux $\langle S^z\rangle$ (the Poynting vector) is given by

$$\frac{\langle h^z\rangle}{\langle S^z\rangle} = \pm\frac{1}{\pi}\frac{e^2}{4\pi\epsilon_0\hbar c}\frac{1}{\hbar\omega}. \tag{10}$$

Hence, the chirality per photon available for the transfer to charged fermions is $\pm\alpha/\pi$, where α is the fine structure constant; the signs refer to the two circular polarizations of light. Of course, the chirality per photon in the beam is ± 1, since photons are massless vector particles. The factor of α/π in equation (10), according to equation (6), describes the coupling of photons to the charged fermions through the chiral anomaly.

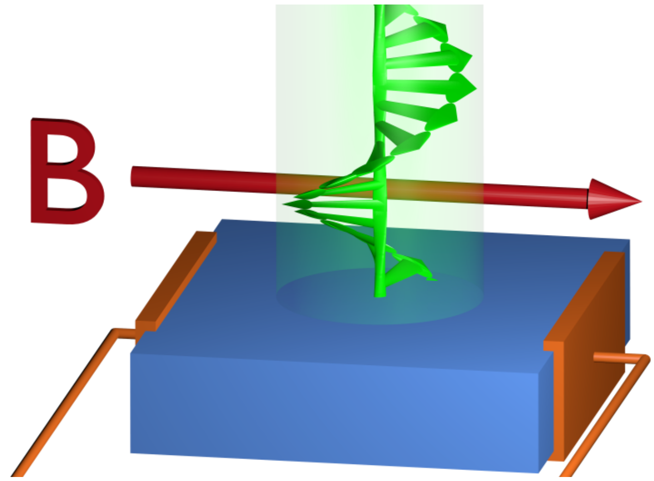

FIG. 1. The CPL incident on a Dirac or Weyl semimetal induces an asymmetry between the number of left- and right-handed chiral quasiparticles. In an external magnetic field, as a consequence of the chiral anomaly, this chiral asymmetry induces a chiral magnetic photocurrent along the direction of the magnetic field.

If CPL is incident on a 3D chiral material, as the light is absorbed by the material, its chirality flux is converted into the chirality of fermions. The total chirality generated per unit area is equal to the chirality flux of light transmitted at the interface. In a real material, if $\omega\tau_{\mathrm{V}} \gg 1$, the chirality will saturate at a constant value proportional to τ_{V} due to chirality relaxation, as dictated by equation (6). Using equation (3), for light incident perpendicular to the interface at $z = 0$,

$$\int_0^\infty \rho_5\,\mathrm{d}z = \tau_{\mathrm{V}}\langle h^z\rangle|_{z=0} = \pm\tau_{\mathrm{V}}\frac{\alpha}{\pi}\frac{I_{\mathrm{in}}}{\hbar\omega}\Re(a_x a_y^*), \tag{11}$$

where I_{in} is the intensity of the incident light and $a_{x,y}$ are the transmission amplitudes of the two linear polarizations. The chiral charge is distributed in the material over a length scale determined by the diffusion length of fermion quasiparticles and the attenuation length of light, but the total chiral charge integrated over the depth is unaffected by this.

The chiral charge density of fermions ρ_5 translates into a chiral chemical potential $\mu_5 \simeq \chi^{-1}\rho_5$, where $\chi = \partial\rho_5/\partial\mu_5$ is the chiral susceptibility. If the whole system is placed in a constant magnetic field $\boldsymbol{B}_{\mathrm{ext}}$ perpendicular to the incident light, a chiral magnetic current[8,9]

$$J_{\mathrm{CME}} = \frac{e^2}{2\pi^2\hbar^2}\boldsymbol{B}_{\mathrm{ext}}\,\mu_5 \tag{12}$$

is generated along the direction of the magnetic field. The linear density κ_{CMP} of the resulting chiral magnetic photocurrent is given by the integral over the depth:

$$\kappa_{\rm CMP} = \int_0^\infty \frac{e^2}{2\pi^2\hbar^2}\,\boldsymbol{B}_{\rm ext}\,\mu_5\,{\rm d}z = \pm\frac{e^2}{2\pi^2\hbar^2}\,B_{\rm ext}\,\frac{\tau_{\rm V}}{\chi}\,\frac{\alpha}{\pi}\,\frac{I_{\rm in}}{\hbar\omega}\,\Re(a_x a_y^*). \tag{13}$$

This formula is the main result of our paper.

For an isotropic chiral material in a weak magnetic field,

$$\chi = \frac{\mu^2}{\hbar^3\pi^2 v_{\rm F}^3} + \frac{k_{\rm B}^2 T^2}{3\hbar^3 v_{\rm F}^3}. \tag{14}$$

For a mid-infrared laser of power 10 mW, intensity $I_{\rm in} = 5\times10^6\,{\rm W/m^2}$, spot diameter $50\,\mu$m, and wavelength $10\,\mu$m incident on an isotropic Dirac semimetal with Fermi velocity $v_{\rm F} = c/500$, chirality relaxation time $\tau_{\rm V} = 10^{-10}\,$s, chemical potential $\mu = 15\,$meV (counted from the Dirac or Weyl node), temperature $T = 77\,$K, and $B = 1\,$T, assuming the light transmission amplitude $|a| = 0.5$, we predict the linear current density in the laser spot to be about 0.5 A/m and the current – about $25\,\mu$A.

Because the chiral magnetic photocurrent will propagate close to the surface of the material, within the light attenuation depth, the bulk of the sample will not be driven. As a result, the magnitude of the observed photocurrent will be reduced by the ratio of the resistances of the sample and the load (leads and contacts). For a sample of resistance $0.1\,\Omega$ (the typical resistivity for known Dirac semimetals is $10^{-5}\,\Omega\,$m) and a load of resistance $10\,\Omega$, the observed current will be suppressed by a factor of 100. The observed photocurrent is thus predicted to be of the order of 200 nA. This can be compared to the photocurrent of about 40 nA observed recently in TaAs[15] that already exceeds, as far as we are aware, the photocurrent observed in other materials currently used for detecting the mid-infrared radiation by a factor of 10-100. In addition, we note that the magnitude of the photocurrent is proportional to ω^{-1} (see equation (13)); therefore, the photocurrent is predicted to be much stronger in the THz frequency range.

The magnitude of the predicted photocurrent is similar to the one predicted recently[23,24] for linearly polarized light in Weyl semimetals with tilted Weyl cones, such as TaAs, in an external magnetic field. The difference between the effect predicted here and in that work is that the chiral magnetic photocurrent considered here should appear only for CPL, but independently of the tilt and the symmetries of Weyl or Dirac cones – the only necessary condition is the existence of chiral fermion quasiparticles. In a material with symmetric cones, for the same intensity, the two circular polarizations will result in currents of equal magnitude but of opposite sign.

The observation of the chiral magnetic photocurrent would provide strong independent evidence for the importance of the chiral anomaly in condensed matter systems. The magnitude of the chiral magnetic photocurrent is predicted to significantly exceed the photocurrents in the currently available devices. This opens possibilities for applications in photonics, optoelectronics, and THz sensing.

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* sahal.kaushik@stonybrook.edu
† evan.philip@stonybrook.edu
‡ dmitri.kharzeev@stonybrook.edu; kharzeev@bnl.gov

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