

Spin-current generation from Coulomb-Rashba interaction in semiconductor bilayers

M. M. Glazov,¹ M. A. Semina,¹ S. M. Badalyan,^{2,3,4} and G. Vignale⁵

¹*Ioffe Physical-Technical Institute of the Russian Academy of Sciences, St. Petersburg, 194021 Russia*

²*School of Physics, Astronomy, and Computational Sciences, George Mason University, Fairfax, Virginia 22030, USA*

³*Department of Physics, University of Antwerp, Groenenborgerlaan 171, B-2020 Antwerpen, Belgium*

⁴*Department of Physics, University of Regensburg, D-93040 Regensburg, Germany*

⁵*Department of Physics and Astronomy, University of Missouri, Columbia, Missouri 65211, USA*

(Received 10 June 2011; published 26 July 2011)

Electrons in double-layer semiconductor heterostructures experience a special type of spin-orbit interaction that arises in each layer from the perpendicular component of the Coulomb electric field created by electron-density fluctuations in the *other* layer. We show that this interaction, acting in combination with the usual spin-orbit interaction, can generate a spin current in one layer when a charge current is driven in the other. This effect is distinct symmetrywise from the spin-Hall drag. The spin current is not, in general, perpendicular to the drive current.

DOI: 10.1103/PhysRevB.84.033305

PACS number(s): 72.10.-d, 71.70.Ej, 72.25.-b, 73.40.-c

I. INTRODUCTION

Spin-orbit interactions in semiconductors are traditionally studied within a one-electron picture, but there are instances in which electron-electron interactions latch onto spin-orbit interactions to produce intriguing effects, which may lead to the creation of radically different spin-based electronic devices.¹⁻³ Of particular interest is the generation of spin currents and spin accumulation by an electric current, through the so-called *spin-Hall effect*.^{4,5} Recently, a special type of spin-Hall effect was predicted to occur in double-layer heterostructures, i.e., two parallel quantum wells separated by an essentially impenetrable potential barrier, with a quasi-two-dimensional electron gas in each layer. This effect, called spin-Hall drag,⁶ consists of the generation of transverse spin accumulation in one layer by an electric current flowing along the other layer and is caused by the component of the Coulomb electric field parallel to the two layers. In this Brief Report, we study a different effect, which is driven by the component of the Coulomb electric field *perpendicular* to the layers. This perpendicular field creates an inhomogeneous Rashba spin-orbit interaction with spatial variation in the plane of the layer.⁷ We refer to this interaction as the Coulomb-Rashba interaction.

The system under study is shown in Fig. 1. A steady electric current is driven in the active layer 2. We show that the interplay of the Coulomb-Rashba interaction with the ordinary cubic Dresselhaus spin-orbit interaction (characteristic of semiconductors of the zinc-blende structure), provides a new mechanism for the generation of a spin current in the passive layer 1. This is particularly remarkable in view of the fact that ordinary spin-Hall effect and spin-Hall drag are *suppressed* by spin precession. But, in this case, the presence of spin precession is absolutely essential to the effect.

The spin-current generation can be understood as a two-stage process, which is schematically described in Fig. 2.¹⁰ In the first stage, the steady current of electrons in layer 2 induces a quadrupolar spin distribution in layer 1. To qualitatively understand how this comes about, observe that Coulomb collisions between electrons in the two layers take

place on average with a positive momentum transfer from layer 2 to layer 1 along the x axis. Moreover, due to the spin-dependent terms in the scattering rate [see Eq. (1) below], the scattering efficiency depends on the relative orientations of electron spin and its initial and final wave vectors, denoted by \mathbf{k} and \mathbf{p} , respectively. Consider electrons with spins parallel or antiparallel to the y axis: $s_y > 0$ and $s_y < 0$, respectively. For $s_y > 0$, the scattering rate is maximal for $k_x + p_x > 0$ and is minimal for $k_x + p_x < 0$. For $s_y < 0$, the situation is reversed: The scattering rate is maximal for $k_x + p_x < 0$ and is minimal for $k_x + p_x > 0$. The strongest transitions are marked by solid arrows in Fig. 2(a). As a result, states with large values of $|k_y|$ ($|k_y| \sim k_F$) become depleted of $s_y > 0$ electrons and filled with $s_y < 0$ electrons. At the same time, states with large values of $|k_x|$ ($|k_x| \sim k_F$) are depleted of $s_y < 0$ electrons and are filled with $s_y > 0$ electrons. Hence, a quadrupolar spin distribution is formed, as shown in Fig. 2(b).

The second stage of the spin-current generation is related to the spin precession in the cubic Dresselhaus field. This is also illustrated in Fig. 2(b), where the green arrows show the direction of the spin precession induced by the cubic Dresselhaus field $\mathbf{\Omega}_k$ [see Eq. (4) below]. The field tilts the spins out of the x - y plane, thus, creating a dipolar distribution of the z component of the spin as shown in Fig. 2(c). As a result, a z -spin current is formed. The characteristic C_{2v} point symmetry of systems with both Rashba and Dresselhaus interactions causes the spin current to be parallel to the driving electric field \mathbf{E} , when \mathbf{E} is along one of the cubic [100] or [010] axes. Notice that this is completely different from the spin-Hall drag current, which is always perpendicular to the electric field. However, the spin current can also be made perpendicular to \mathbf{E} , by orienting \mathbf{E} is along one of the principal axes [110] or $[\bar{1}10]$.

In the rest of the Brief Report, we present some details of our theoretical analysis and provide a numerical estimate for the size of the effect. The salient conclusions are as follows: (i) The spin transresistivity is proportional to $(T/E_F)^2$, which is characteristic of Coulomb drag phenomena,^{6,12,13} (ii) it is inversely proportional to the fourth power of the interlayer separation, and (iii) it is parametrically stronger than the spin-Hall drag⁶ in the clean limit since the transresistivity

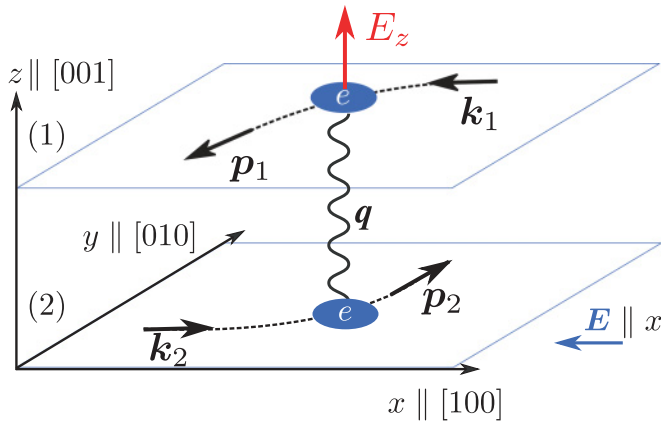


FIG. 1. (Color online) Scheme of the device under study: An electric field \mathbf{E} is applied to the active layer (2), and a spin current is generated in the passive layer (1). Blue circles depict electrons, and the wavy line shows the interlayer Coulomb interaction. Arrows marked $\mathbf{k}_1, \mathbf{k}_2$ and $\mathbf{p}_1, \mathbf{p}_2$ are the wave vectors in the initial and final states. The vertical arrow emphasizes the relevant component of the interlayer Coulomb field.

is proportional to $\Omega\tau$, whereas, the T^2 contribution to the spin-Hall drag is independent of τ .

II. THEORY

The part of the spin-orbit-dressed Coulomb interaction,¹⁴ which is relevant to the effect described in this Brief Report, has matrix elements,^{8,15,16}

$$M(\mathbf{k}_1 s_1, \mathbf{k}_2 s_2 \rightarrow \mathbf{p}_1 s'_1, \mathbf{p}_2 s'_2) = \left\{ V_q - \frac{\lambda^2}{2} V_q (q + q_s) \langle \chi_{s_1} | \chi_{s'_1} \rangle [\hat{\sigma}_1 \times (\mathbf{p}_1 + \mathbf{k}_1)]_z - [\hat{\sigma}_2 \times (\mathbf{p}_2 + \mathbf{k}_2)]_z \langle \chi_{s_2} | \chi_{s'_2} \rangle \right\} \delta_{\mathbf{k}_1 + \mathbf{k}_2, \mathbf{p}_1 + \mathbf{p}_2}, \quad (1)$$

where subscripts 1 and 2 refer to electrons in layers 1 and 2, respectively, $\hat{\sigma}_1$ and $\hat{\sigma}_2$ are the spin operators of the first and second electron acting on spinors $|\chi_{s_1}\rangle, |\chi_{s'_1}\rangle$ and $|\chi_{s_2}\rangle, |\chi_{s'_2}\rangle$, respectively. The normalization area is set to unity. Here, $\mathbf{q} = \mathbf{p} - \mathbf{k}$ is the transferred wave vector,

$$V_q = \frac{2\pi e^2}{\kappa} \frac{q e^{-qL}}{(q + q_s)^2 - q_s^2 e^{-2qL}}$$

is the Fourier transform of the screened Coulomb interaction between the layers, L is the distance between the layers, κ is a background dielectric constant, and q_s is the screening wave vector. The strength of the spin-orbit interaction is controlled by the effective Compton wavelength λ for the semiconductor,

$$\lambda^2 = -\frac{P^2}{3} \frac{\Delta(2E_g + \Delta)}{E_g^2(E_g + \Delta)^2}, \quad (2)$$

where P is the Kane parameter¹⁷ and E_g and Δ are the band-gap and the spin-orbit splitting of the valence band. In GaAs, $\lambda^2 \approx 5 \text{ \AA}^2$. Notice that, in writing Eq. (1), we have taken only the Rashba-like contribution arising from the component of the Coulomb field perpendicular to the layers into account. The additional term coming from the component of the Coulomb

field parallel to the layers was discussed in Refs. 6,18 and 19 and is not shown in Eq. (1).

As discussed in Sec. I and in the caption of Fig. 2, a steady current driven in the active layer (2) produces, via the spin-orbit dressed Coulomb interaction, a quadrupolar distribution of spin in the passive layer. The spin-generation rate in the passive layer can be calculated by means of the Fermi golden rule with the matrix element of interlayer electron-electron interaction given by Eq. (1).^{20,21} For first order in λ^2 , this gives

$$\mathbf{g}_k = \frac{4\pi\hbar e\tau}{mk_B T} \sum_{k'p'p} \delta_{\mathbf{k}+\mathbf{k}', \mathbf{p}+\mathbf{p}'} \delta(\mathcal{E}_k + \mathcal{E}_{k'} - \mathcal{E}_p - \mathcal{E}_{p'}) \times U_{k-p} [(\mathbf{p} + \mathbf{k}) \times \hat{\mathbf{z}}] [\mathbf{E} \cdot (\mathbf{p} - \mathbf{k})] \times f_k f_{k'} (1 - f_p)(1 - f_{p'}), \quad (3)$$

where $U_q \equiv \lambda^2 (q + q_s) |V_q|^2$, \mathbf{E} is the electric field acting on the carriers in the active layer, m is the effective electron mass, $k_B T$ is the temperature measured in the energy units, $\mathcal{E}_k = \hbar^2 k^2 / (2m)$ is the electron dispersion, f_k is the Fermi-Dirac distribution function, and $\hat{\mathbf{z}}$ is the unit vector normal to the layers. We have also assumed that equilibrium densities, effective masses, and Fermi energies are the same in the two layers. In deriving Eq. (3), we took into account only the linear-in- \mathbf{E} nonequilibrium correction to the distribution function. It is easy to show that, due to the presence of the two factors $\mathbf{E} \cdot (\mathbf{p} - \mathbf{k})$ and $(\mathbf{p} + \mathbf{k}) \times \hat{\mathbf{z}}$, the angular dependence of \mathbf{g}_k is

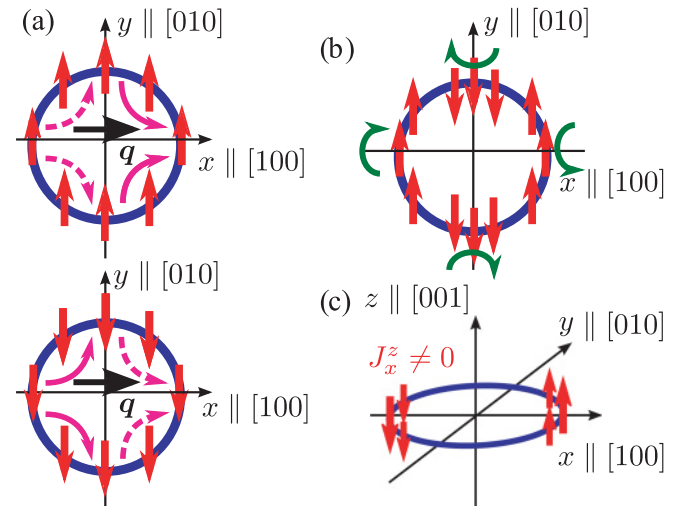


FIG. 2. (Color online) Schematic of the spin-current generation in the passive layer.¹⁰ The circle is the Fermi surface, and the arrows are the electron spins. Panel (a) shows the scattering stage of the process: The top part shows electrons with spin parallel to the y axis, and the bottom part shows electrons with spin antiparallel to the y axis. The arrows show the spin-dependent scattering process: Solid arrows indicate the stronger transitions, having $k_x + p_x > 0$ for $s_y > 0$ and $k_x + p_x < 0$ for $s_y < 0$, while dashed arrows indicate the weaker transitions. Scattering processes that increase the x component of the wave vector (i.e., with $q_x > 0$) dominate due to the current flowing in the active layer. (b) Precession stage of the process. The resulting spin distribution after the scattering contains second angular harmonics. Green arrows demonstrate the spin precession caused by the Dresselhaus field Ω_k , Eq. (4). Electron spins with opposite wave vectors precess in opposite directions. The resulting dipolar distribution of electron spins is presented in panel (c).

an angular harmonic of order 2, i.e., we have $g_{x,k} \propto \sin 2\varphi_k$ and $g_{y,k} \propto -\cos 2\varphi_k$, where φ_k is the angle of \mathbf{k} with the x axis: This is the quadrupolar pattern of spin generation mentioned in Sec. I and described in Fig. 2.

The spin dynamics in the passive layer is governed by the spin-orbit splitting of the energy spectrum. Since the first-harmonic components of the spin splitting (arising from linear-in- \mathbf{k} terms) do not result in a dc spin Hall current,^{11,22,23} we only take into account the third angular harmonics of the k^3 Dresselhaus term, which is inevitable in any zinc-blende structure. In cubic axes with $x \parallel [100]$ and $y \parallel [010]$, these take the form

$$\Omega_{x,k} = -\Omega_3 \cos 3\varphi_k, \quad \Omega_{y,k} = -\Omega_3 \sin 3\varphi_k, \quad (4)$$

where $\Omega_3 = \gamma_c k^3 / (2\hbar)$, γ_c is the bulk Dresselhaus constant. The electron-spin distribution function s_k is determined by a kinetic equation, which, in the steady state, takes the form¹¹

$$s_k \times \Omega_k + \frac{s_k}{\tau} = \mathbf{g}_k. \quad (5)$$

Here, we have assumed that all spin-independent scattering processes can be characterized by a single relaxation time τ . The solution of Eq. (5) is

$$s_{z,k} = \frac{(\Omega_k \times \mathbf{g}_k)_z}{\Omega_k^2 + 1/\tau^2}. \quad (6)$$

We note that such a simple form of the solution results from keeping only the third angular harmonics in Ω_k . The solution in the general case can be constructed following Refs. 11 and 24: In such a case, the overall spin will be smaller due to faster spin relaxation caused by the linear-in- \mathbf{k} Dresselhaus term, and the spin distribution may have a more complicated form due to the anisotropy of the spin splitting.

By definition, the current density of the spin- z component is

$$\mathbf{J}^z = \sum_k s_{z,k} \mathbf{v}_k, \quad (7)$$

where $\mathbf{v}_k = \hbar \mathbf{k} / m$ is the electron velocity. Substituting the expression for \mathbf{g}_k , obtained from Eq. (3), into Eq. (6) for $s_{z,k}$, and then $s_{z,k}$ in Eq. (7), we arrive at our main result,

$$\begin{aligned} \mathbf{J}^z = & -\frac{2\pi\hbar e\tau^3}{mk_B T} \sum_{kk'pp'} \mathbf{v}_k \delta_{k+k', p+p'} \delta(\mathcal{E}_k + \mathcal{E}_{k'} - \mathcal{E}_p - \mathcal{E}_{p'}) \\ & \times 2U_{k-p} [\Omega_k \cdot (\mathbf{p} + \mathbf{k})][\mathbf{E} \cdot (\mathbf{p} - \mathbf{k})] \\ & \times f_k f_{k'} (1 - f_p)(1 - f_{p'}). \end{aligned} \quad (8)$$

Equation (8) can be recast in the standard form

$$J_i^z = G_{ij} E_j, \quad (9)$$

where G_{ij} is the spin-drag conductivity. Our system is characterized by the C_{2v} point symmetry group because we take (i) the Dresselhaus field and (ii) the Rashba-like interaction associated with the perpendicular-to-plane component of the Coulomb field into account. On the basis of the principal axes $x_1 \parallel [1\bar{1}0]$ and $y_1 \parallel [110]$, Eq. (9) can be written in terms of two independent constants G and G_1 as

$$J_{x_1}^z = (G + G_1)E_{y_1}, \quad J_{y_1}^z = (G - G_1)E_{x_1}. \quad (10)$$

Equation (10) gives the full phenomenological picture of the spin drag in the presence of the spin-orbit interaction. The drag mechanism described here produces $G \neq 0$ and $G_1 = 0$. This rather peculiar situation implies that the spin current flows parallel to the electric field if the latter is applied along the $[100]$ or $[010]$ directions, but, in general, it has both a parallel and a perpendicular component. It is only when the electric field is along one of the principal axes that we get a pure transverse current, but even in this case, it is in sharp contrast with the spin-Hall drag,⁶ namely, an electric field along x_1 produces a spin current along y_1 , but an electric field along y_1 produces a spin current along x_1 , rather than $-x_1$. These unique symmetry-related features will help distinguish the predicted different effect.

We now come to the quantitative evaluation of G . The expression (8) for the spin-current density is similar to the expression encountered in the calculation of the ordinary drag current.¹³ The main difference is that the integrand depends not only on the momentum transfer $\mathbf{q} = \mathbf{p} - \mathbf{k}$, but also on the sum of the initial and final momenta $\mathbf{p} + \mathbf{k}$. For simplicity, here, we focus on the case of well-separated layers $k_F L \gg 1$, where the momentum transfer is typically small ($q \sim L^{-1} \ll k_F$), and, therefore, one can replace $\mathbf{k} + \mathbf{p} \approx 2\mathbf{k}_F$, producing an error of order q/k_F . With this approximation, the standard method of evaluation can be applied, and, after converting the wave-vector sums into integrals, and for low enough temperatures $k_B T \ll E_F$, we obtain

$$G = -\frac{\hbar^3 e \Omega_3 k_F \tau^3}{2\pi m^2 k_B T} \int \frac{d\mathbf{q}}{(2\pi)^2} \int_{-\infty}^{+\infty} d\omega \frac{\text{Im}[\chi_0(q, \omega)]^2}{\sinh^2(\frac{\hbar\omega}{2k_B T})} U_q q^2, \quad (11)$$

where $\chi_0(q, \omega)$ is the noninteracting density-density response function of the two-dimensional electron gas.²⁵ Integrating over frequency, we get

$$G = -\frac{\lambda^2 \Omega_3 e^5 k_F^4 \tau^3}{4\pi^2 \kappa^2 \hbar^2} \left(\frac{k_B T}{E_F}\right)^2 I_2(q_s L), \quad (12)$$

with

$$I_2(x) = \frac{2\pi^2}{3k_F L} \int_0^\infty \frac{y^3(y+x)e^{-2y}}{[(x+y)^2 - x^2 e^{-2y}]^2} dy, \quad (13)$$

for $k_F L \gg 1$. For $q_s L \gg 1$, the integral Eq. (13) is evaluated to be $I_2(q_s L) \approx 2.96 / (k_F q_s^3 L^4)$, as a result, $G \propto 1/L^4$, as in the ordinary drag effect.^{12,13}

III. RESULTS AND DISCUSSION

It is instructive to estimate the drag resistivity that controls the observable spin accumulation. This is determined as follows. First, we express the external field E via the current density generated by this field in the active layer $j = \sigma E$, where $\sigma = ne^2\tau/m$ is the Drude conductivity of the electrons and n is their density. Then, we observe that the spin current is associated with some effective electric spin field E_s ,

$$\frac{2e}{\hbar} J^z = \sigma E_s.$$

As a result, we obtain

$$E_s = \frac{2eJ^z}{\hbar\sigma} = \frac{2eG}{\hbar\sigma^2}j = \rho_s j.$$

Hence, the drag resistivity is given by

$$\rho_s = -\frac{2\hbar}{e^2}(\lambda k_F)^2 \Omega_3 \tau \left(\frac{k_B T}{E_F}\right)^2 \left(\frac{e^2}{\kappa \hbar v_F}\right)^2 I_2(q_s L), \quad (14)$$

where v_F is the Fermi velocity. Notice that, at variance with the (side-jump) spin-Hall drag resistivity, the present result is proportional to the momentum relaxation time τ , and, therefore, it is parametrically dominant in the clean limit $\tau \rightarrow \infty$.

For the following values of the parameters: electron density $n = 2 \times 10^{11} \text{ cm}^{-2}$, $\tau = 40 \text{ ps}$ (which corresponds to a mobility $\mu = 10^6 \text{ cm}^2 \text{ V}^{-1} \text{ s}^{-1}$ in a GaAs quantum well with $m = 0.067m_0$), bulk Dresselhaus splitting constant $\gamma_c = 20 \text{ eV \AA}^3$, $\lambda^2 = 5 \text{ \AA}^2$, and $\kappa = 13$, we obtain $(\lambda k_F)^2 \approx 6 \times 10^{-4}$, $\Omega_3 \tau \approx 0.9$, and $e^2/\kappa \hbar v_F \approx 1$. Thus, taking $I_2 \sim 1$, we have $\rho_s \sim 4\left(\frac{k_B T}{E_F}\right)^2 \text{ Ohm}$, which for $(k_B T/E_F) = 0.1$ is 0.04 Ohm , i.e., about two times larger than the conservative estimate for the spin-Hall drag resistivity from the side jump given in Ref. 6. In

InAs-based structures, the spin-orbit coupling parameter λ^2 is about an order of magnitude higher than in GaAs, and Ω_3 is also larger. This makes narrow-band semiconductor bilayers particularly suitable for the observation of the spin-current injection.

In conclusion, we have described a coupling mechanism, partly Coulomb and partly spin orbit, through which a spin current can be injected, or a spin accumulation can be induced, in an electron layer by a regular electric current flowing in an adjacent layer. The coupling can play a role in the design of circuits in which an electric current must be converted into a spin current and vice versa.

ACKNOWLEDGMENTS

M.M.G. and M.A.S. are grateful to RFBR, EU projects Spinoptronics and POLAPHEN, and the ‘‘Dynasty’’ Foundation—ICFPM for financial support. S.M.B. acknowledges support from EU Grant No. PIIF-GA-2009-235394, the DFG SFB 689, and the Belgium Science Policy (IAP). G.V. acknowledges support from NSF Grant No. DMR-0705460.

¹D. Awschalom and N. Samarth, *Phys.* **2**, 50 (2009).

²M. I. Dyakonov, ed. *Spin Physics in Semiconductors* (Springer-Verlag, Berlin, 2008).

³J. Fabian, A. Matos-Abiague, C. Ertler, P. Stano, and I. Žutić, *Acta Phys. Slov.* **57**, 565 (2007).

⁴M. I. Dyakonov and V. I. Perel’, *Phys. Lett.* **35A**, 459 (1971).

⁵Y. K. Kato, R. C. Myers, A. C. Gossard, and D. D. Awschalom, *Science* **306**, 1910 (2004); J. Wunderlich, B. Kaestner, J. Sinova, and T. Jungwirth, *Phys. Rev. Lett.* **94**, 047204 (2005).

⁶S. M. Badalyan and G. Vignale, *Phys. Rev. Lett.* **103**, 196601 (2009).

⁷We note that an inhomogeneous Rashba interaction also arises from impurities in the doping layer and was previously shown^{8,9} to be part of the electron-impurity spin-orbit coupling.

⁸M. M. Glazov, E. Ya. Sherman, and V. K. Dugaev, *Physica E (Amsterdam)* **42**, 2157 (2010).

⁹N. S. Averkiev, L. E. Golub, and M. Willander, *J. Phys.: Condens. Matter* **14**, R271 (2002).

¹⁰In monolayers, similar arguments for electron-impurity and electron-phonon scatterings were presented in Ref. 11. Spin current generation solely due to inhomogeneous Rashba interaction in quantum wells was studied in Ref. 8.

¹¹S. Tarasenko, *JETP Lett.* **84**, 199 (2006).

¹²T. J. Gramila, J. P. Eisenstein, A. H. MacDonald, L. N. Pfeiffer, and K. W. West, *Phys. Rev. Lett.* **66**, 1216 (1991).

¹³A.-P. Jauho and H. Smith, *Phys. Rev. B* **47**, 4420 (1993).

¹⁴See Supplemental Material at <http://link.aps.org/supplemental/10.1103/PhysRevB.84.033305> for a detailed derivation.

¹⁵M. M. Glazov and V. D. Kulakovskii, *Phys. Rev. B* **79**, 195305 (2009).

¹⁶M. M. Glazov, M. A. Semina, and E. Y. Sherman, *Phys. Rev. B* **81**, 115332 (2010).

¹⁷R. Winkler, *Spin-Orbit Effects in Two-Dimensional Electron and Hole Systems* (Springer, New York, 2003).

¹⁸P. Boguslawski, *Solid State Commun.* **33**, 389 (1980).

¹⁹M. M. Glazov, *J. Phys.: Condens. Matter* **22**, 025301 (2010).

²⁰I. D’Amico and G. Vignale, *Phys. Rev. B* **65**, 085109 (2002).

²¹M. M. Glazov and E. L. Ivchenko, *JETP* **99**, 1279 (2004); *Europhys. Lett.* **87**, 57005 (2009).

²²O. Chalaev and D. Loss, *Phys. Rev. B* **71**, 245318 (2005).

²³O. V. Dimitrova, *Phys. Rev. B* **71**, 245327 (2005).

²⁴M. M. Glazov and L. E. Golub, *Phys. Rev. B* **77**, 165341 (2008).

²⁵G. F. Giuliani and G. Vignale, *Quantum Theory of Electron Liquid* (Cambridge University Press, Cambridge, UK, 2005).