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Bernard

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# Giant Collective Spin-Orbit Field in a Quantum Well: Fine Structure of Spin Plasmons

F. Baboux,<sup>1,\*</sup> F. Perez,<sup>1</sup> C. A. Ullrich,<sup>2</sup> I. D'Amico,<sup>3</sup> J. Gómez,<sup>1,†</sup> and M. Bernard<sup>1</sup>

<sup>1</sup>*Institut des Nanosciences de Paris, CNRS/Université Paris VI, Paris 75005, France*

<sup>2</sup>*Department of Physics and Astronomy, University of Missouri, Columbia, Missouri 65211, USA*

<sup>3</sup>*Department of Physics, University of York, York YO10 5DD, United Kingdom*

We employ inelastic light scattering with magnetic fields to study intersubband spin plasmons in a quantum well. We demonstrate the existence of a giant collective spin-orbit (SO) field that splits the spin-plasmon spectrum into a triplet. The effect is remarkable as each individual electron would be expected to precess in its own momentum dependent SO field, leading to D'yakonov-Perel' dephasing. Instead, many-body effects lead to a striking organization of the SO fields at the collective level. The macroscopic spin moment is quantized by a uniform collective SO field, 5 times higher than the individual SO field. We provide a momentum-space cartography of this field.

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Spin-orbit (SO) coupling arises from relativity: the spin of an electron moving at a velocity  $\mathbf{v}$ , in a static electric field  $\mathbf{E}$ , sees a magnetic field  $\mathbf{B}_{\text{SO}} = -\frac{1}{c^2}\mathbf{v} \times \mathbf{E}$  ( $c$  is the speed of light) [1]. This magnetic field splits the energy levels of atoms, giving rise to their fine structure [2]. For an ensemble of itinerant electrons in solids, such a simple quantizing effect cannot be expected, because of the distribution of velocities: Momentum-dependent SO fields cause each individual electronic spin to precess with its own axis, which destroys spin coherence (D'yakonov-Perel' [DP] decoherence [3]). This sets practical limitations on many proposed applications in emerging quantum technologies such as spintronics [4–8].

However, this DP picture is appropriate only for situations where the macroscopic spin is carried by individual electrons, which is often the case [5, 8–11]. Here, we demonstrate that Coulomb interaction, playing a central role in *collective* spin excitations, can drastically modify this picture, and give rise to macroscopic quantum objects. We will focus on intersubband spin plasmons in doped semiconductor quantum wells, which, as we shall see, are ideal to study the interplay of SO coupling and Coulomb interactions.

In a III-V quantum well, internal SO fields arise from the lack of an inversion center of the crystalline unit cell, and from an asymmetric confining potential [12], referred to as Dresselhaus [13] and Rashba [14] fields, respectively. Hence, a conduction electron with momentum  $\mathbf{k}$ , moving in the plane of a [001]-oriented quantum well, experiences a SO magnetic field

$$\mathbf{B}_{\text{SO}}(\mathbf{k}) = \frac{2\alpha}{g\mu_B} \begin{pmatrix} k_y \\ -k_x \end{pmatrix} + \frac{2\beta}{g\mu_B} \begin{pmatrix} k_x \\ -k_y \end{pmatrix} \quad (1)$$

(to lowest order in  $\mathbf{k}$ ), for coordinate systems with  $\hat{x} \parallel [100]$  and  $\hat{y} \parallel [010]$ . Here,  $\alpha$  and  $\beta$  are the Rashba and linear Dresselhaus coupling constants [12], respectively,  $g$  is the electron g-factor, and  $\mu_B$  the Bohr magneton.  $\mathbf{B}_{\text{SO}}$  produces an intrinsic  $\mathbf{k}$ -dependent spin splitting [9,

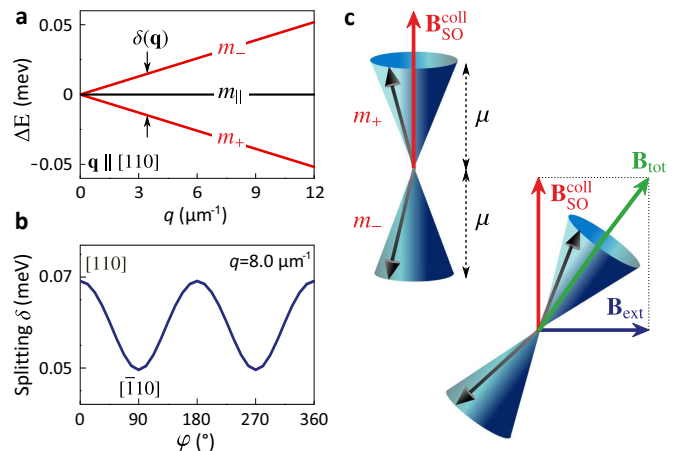


FIG. 1. Fine structure model of ISB spin plasmons. (a) Three-fold splitting, induced by SO coupling, of the ISB spin plasmon modes.  $\Delta E$  denotes the difference of the mode energies with and without SO coupling, calculated for the studied GaAs quantum well, and  $q$  is the magnitude of the plasmon momentum (here  $\mathbf{q} \parallel [110]$ ). The splitting  $\delta$  between the transverse  $m_{\pm}$  modes is almost linear in  $q$ . (b) For a fixed  $q = 8.0 \mu\text{m}^{-1}$ , calculated modulation of the splitting  $\delta$  with the in-plane orientation of  $\mathbf{q}$ , labeled by the angle  $\varphi$  to [110]. (c) Sketch of the proposed interpretation of the transverse ISB spin plasmons  $m_{\pm}$ , as the precession of antiparallel  $\mu$  collective magnetic moments about  $\mathbf{B}_{\text{SO}}^{\text{coll}}$  at zero external field (left), and about the superposition  $\mathbf{B}_{\text{SO}}^{\text{coll}} + \mathbf{B}_{\text{ext}}$  when an external magnetic field  $\mathbf{B}_{\text{ext}}$  is applied (right).

10] and a  $\mathbf{k}$ -dependent spin orientation [6, 15] of single-electron conduction states.

In such a system, electrons can exhibit collective spin dynamics when excited from the first to the second subband of the quantum well. These so-called intersubband (ISB) spin plasmons, which arise from Coulomb interactions, are energetically well separated from the continuum of ISB single-particle excitations [16, 17]. In the absence of a transferred momentum  $\mathbf{q}$  and external magnetic field  $\mathbf{B}_{\text{ext}}$ , time reversal symmetry, together with

the [001]-axis symmetry of the quantum well, average out the  $\mathbf{k}$ -dependent  $\mathbf{B}_{\text{SO}}$ . Hence, no macroscopic SO force is acting on the electron gas, and the spin plasmons are degenerate. However, when transferring an in-plane momentum  $\mathbf{q}$  to the electron gas, the translation symmetry is broken and  $\mathbf{B}_{\text{SO}}(\mathbf{k})$  does not average out anymore.

In this situation, it has been predicted [18, 19] that despite the spread of  $\mathbf{B}_{\text{SO}}(\mathbf{k})$ , a collective SO magnetic field  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$  emerges, splitting the spin plasmon branch into three modes [Fig. 1(a)]: one longitudinal oscillation mode ( $m_{\parallel}$ ) and two transverse precession modes ( $m_{+}$  and  $m_{-}$ ). In the present work we focus on these transverse modes, whose frequencies are shifted in opposite directions by SO coupling. We propose that due to Coulomb interaction, these modes behave as *macroscopic quantum objects*, characterized by a collective spin magnetic moment  $\mathbf{M}$ , and thus subject to an interaction energy  $W(\mathbf{q}) = -\mathbf{M} \cdot \mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$ . Within this framework, the downward (upward) energy shift of the  $m_{+}$  ( $m_{-}$ ) mode is explained by its projected magnetic moment being parallel (antiparallel) to the quantizing field  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$  [Fig. 1(c), left]. Then, in the presence of an external magnetic field, we expect both fields to superpose [5, 11] [Fig. 1(c), right] and the interaction energy to become

$$W(\mathbf{q}) = -\mathbf{M} \cdot (\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q}) + \mathbf{B}_{\text{ext}}). \quad (2)$$

We will demonstrate that this fine structure model correctly captures the physics of the ISB spin plasmons.

We carry out inelastic light scattering (ILS) measurements in a [001]-oriented, asymmetrically modulated GaAs/AlGaAs quantum well. The electron density is  $2.3 \times 10^{11} \text{ cm}^{-2}$ , and the mobility  $2 \times 10^7 \text{ cm}^2 \text{ V}^{-1} \text{ s}^{-1}$  at the working temperature  $T \simeq 2 \text{ K}$  (superfluid Helium) [20]. ILS [16, 21] is a powerful tool to study spin excitations at a given transferred momentum  $\mathbf{q}$  (Fig. 2(a), inset) [20]. Standard selection rules [16] allow us to address the various types of inter-subband excitations individually. As shown in the spectra of Fig. 2(a) (top), the charge plasmon is observed only when the incident and scattered photon have parallel polarizations (polarized spectra), while the spin plasmon appears when they have orthogonal polarizations (depolarized spectra). The single-particle excitations continuum appears in both configurations (here as a shoulder of the charge plasmon peak). From now on we focus on the spin plasmon peak, obtained in the depolarized geometry where only the transverse modes,  $m_{+}$  and  $m_{-}$  are probed. Typical spectra, taken in the absence of an external magnetic field, are presented in Fig. 2(a) (bottom). These are obtained for a momentum of fixed magnitude  $q = 8.0 \mu\text{m}^{-1}$ , but various in-plane orientations, labeled by the angle  $\varphi$  between  $\mathbf{q}$  and the [110] direction of the quantum well. The spectra exhibit a single, quasi-Lorentzian peak, of full width at half-maximum (FWHM)  $w$ ; when plotting  $w$  for various  $\varphi$  [Fig. 2(c)],  $w$  is modulated quasi-sinusoidally with a period  $\pi$ . This

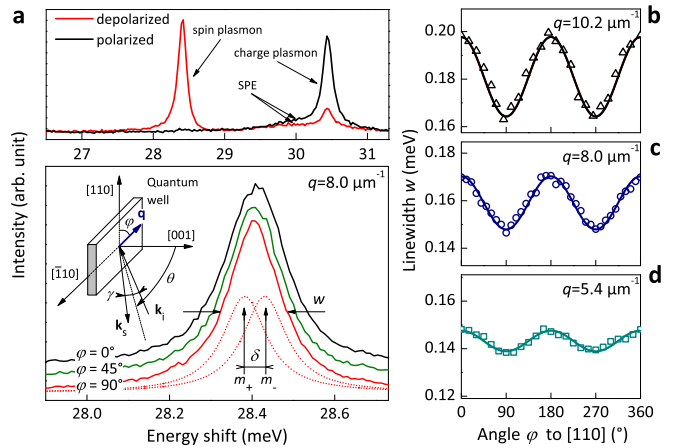


FIG. 2. Anisotropic splitting of the ISB spin plasmon modes. (a) Top panel: Inelastic light scattering spectrum of the ISB excitations, in polarized (black) and depolarized (red) geometry. Bottom panel: Depolarized spectra obtained at fixed  $q = 8.0 \mu\text{m}^{-1}$ , by varying the in-plane angle  $\varphi$  measured from [110] (vertical offset for clarity). The single, quasi-Lorentzian peak observed is the sum of two Lorentzians (red dashed lines for the  $\varphi = 90^\circ$  spectrum) of same amplitude and linewidth, corresponding to the transverse spin plasmons modes  $m_{+}$  and  $m_{-}$  split by an amount  $\delta$ . Inset: scattering geometry showing angle definitions;  $\mathbf{k}_i$  and  $\mathbf{k}_s$  are the incoming and scattered light wavevectors. (b-d) Variation of the linewidth  $w$  with  $\varphi$  for  $q = 10.2, 8.0$  and  $5.4 \mu\text{m}^{-1}$  respectively. Lines: theory (see text).

modulation is characteristic of the two-fold symmetry of the SO splitting [Fig. 1(b)], with a maximum along [110] ( $\varphi = 0^\circ$ ) and a minimum along  $\bar{[110]}$  ( $\varphi = 90^\circ$ ). Furthermore, as seen in Fig. 2(b)-(d), the amplitude of the modulation decreases with decreasing  $q$ , in agreement with Fig. 1(a). Both characteristics confirm the SO origin of the modulation. This suggests that the observed Raman line is the sum of two Lorentzian peaks corresponding to the transverse spin plasmon modes  $m_{+}$  and  $m_{-}$ , split by  $\delta$  (Fig. 2(a), red dashed lines). By independently determining the FWHM of the latter peaks, we will extract the splitting  $\delta(\mathbf{q})$  by deconvolution, and demonstrate the consistency of our model.

We determine the collective SO field  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$  by applying an external magnetic field  $\mathbf{B}_{\text{ext}}$ . Since  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$  is expected [19] to lie in the plane of the quantum well for our [001]-oriented sample,  $\mathbf{B}_{\text{ext}}$  will be applied in the well plane (quasi-Voigt geometry). If the spin plasmon moment  $\mathbf{M}$  is of quantum nature, its energy levels will be quantized by the total field  $\mathbf{B}_{\text{tot}} = \mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q}) + \mathbf{B}_{\text{ext}}$ . Following Eq. (2), the splitting  $\delta$  will then be given by

$$\delta = 2\mu B_{\text{tot}} = 2\mu \sqrt{(B_{\text{ext}} + \mathbf{B}_{\text{SO}}^{\text{coll}} \cdot \mathbf{u})^2 + (\mathbf{B}_{\text{SO}}^{\text{coll}} \times \mathbf{u})^2}, \quad (3)$$

where  $\mu$  is the *quantized* value of the spin plasmon magnetic moment and  $\mathbf{u}$  is a unit vector parallel to the direction of  $\mathbf{B}_{\text{ext}}$ .

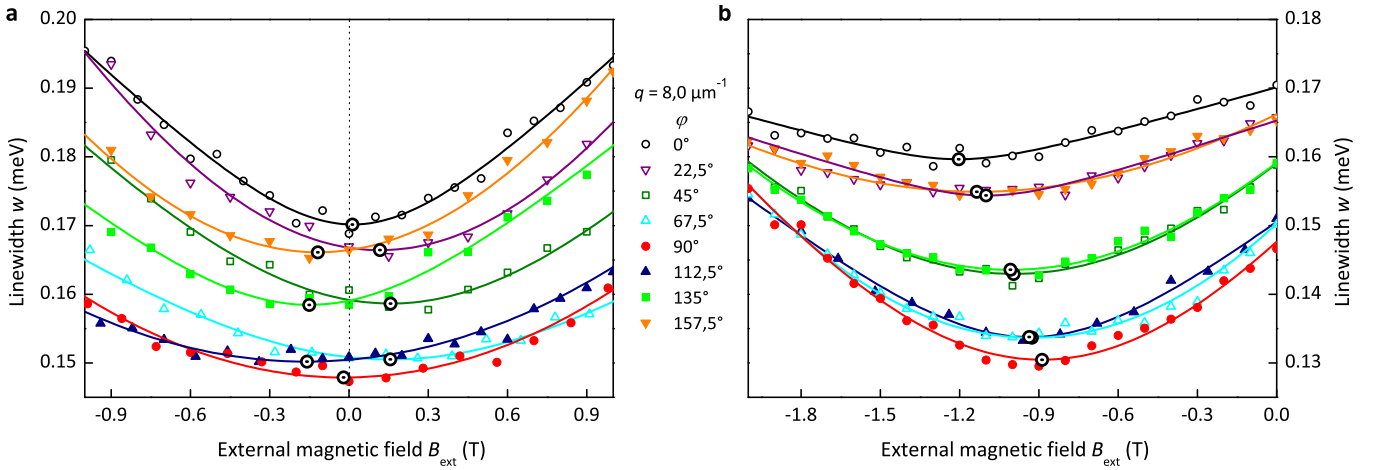


FIG. 3. Variation of the composite linewidth  $w$  with the external magnetic field  $\mathbf{B}_{\text{ext}}$ . (a)  $w(B_{\text{ext}})$  plots obtained in the configuration  $\mathbf{B}_{\text{ext}} \parallel \mathbf{q}$  for a fixed  $q = 8.0 \mu\text{m}^{-1}$  and various in-plane angles  $\varphi$  (measured from  $[110]$ ). (b) Corresponding  $w(B_{\text{ext}})$  plots obtained for  $\mathbf{B}_{\text{ext}} \perp \mathbf{q}$ . Lines are guides for the eyes. Each  $w(B_{\text{ext}})$  plot is symmetric about a certain value of the external field (marked by a dotted circle) which cancels the corresponding component of the collective SO field  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$ .

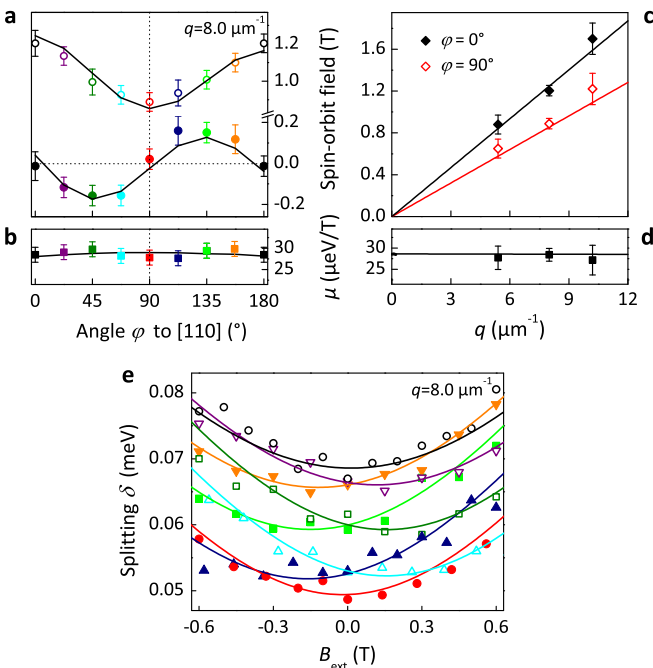


FIG. 4. SO collective field and magnetic moment. (a) Components of the collective SO field parallel ( $B_{\text{SO},\parallel}^{\text{coll}}$ , filled circles) and perpendicular ( $B_{\text{SO},\perp}^{\text{coll}}$ , open circles) to  $\mathbf{q}$  for  $q = 8.0 \mu\text{m}^{-1}$ , as extracted from the data of Fig. 3, and compared with theory (lines). (b) Spin plasmon magnetic moment  $\mu$ , experimental (squares) and theoretical (line). (c) Minimum (open diamonds) and maximum (filled diamonds) SO field  $|\mathbf{B}_{\text{SO}}^{\text{coll}}|$  versus  $q$ , compared to theoretical values (lines). (d) Spin plasmon magnetic moment averaged over  $\varphi$ , experimental (squares) and theoretical (line), as a function of  $q$ . (e) Variation of the SO splitting  $\delta$  with external magnetic field  $\mathbf{B}_{\text{ext}} \parallel \mathbf{q}$ , for  $q = 8.0 \mu\text{m}^{-1}$  (symbols, same as Fig. 3). The experimental data are very well reproduced by Eq. (3) (lines).

For a given  $\mathbf{q}$ , we record a series of spectra at varying  $B_{\text{ext}}$ , with  $\mathbf{B}_{\text{ext}}$  applied successively along two crossed directions:  $\mathbf{B}_{\text{ext}} \perp \mathbf{q}$  and  $\mathbf{B}_{\text{ext}} \parallel \mathbf{q}$ . Figures 3(a) and 3(b), respectively, present the *composite* linewidth  $w$  as a function of  $B_{\text{ext}}$  for these two configurations. The various plots are obtained for fixed  $q = 8.0 \mu\text{m}^{-1}$ , and a set of 8 angles  $\varphi$ , spaced by  $22.5^\circ$  within a period  $\pi$ . Each plot exhibits a clear minimum for a certain value of  $B_{\text{ext}}$ , and is symmetric with respect to that minimum.

According to Eq. (3), each minimum corresponds to the situation where  $\mathbf{B}_{\text{ext}}$  exactly cancels the component of  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$  parallel to it,  $B_{\text{ext}} = -\mathbf{B}_{\text{SO}}^{\text{coll}} \cdot \mathbf{u}$ . Using this criterion, we extract the component  $B_{\text{SO},\parallel}^{\text{coll}}$  of the collective SO field parallel to  $\mathbf{q}$  from the plots of Fig. 3(a), and the perpendicular component  $B_{\text{SO},\perp}^{\text{coll}}$  from the plots of Fig. 3(b). Figure 4(a) presents the values for  $B_{\text{SO},\parallel}^{\text{coll}}$  (filled circles) and  $B_{\text{SO},\perp}^{\text{coll}}$  (open circles). We find that  $B_{\text{SO},\parallel}^{\text{coll}}$  is antisymmetric about the  $[\bar{1}10]$  direction ( $\varphi = 90^\circ$ ), and  $B_{\text{SO},\perp}^{\text{coll}}$  is symmetric.

To push the analysis further, we need experimental access to the SO splitting  $\delta$ . This can be done by determining the linewidth of the  $m_{\pm}$  modes [see Fig. 2(a)]. The latter is inferred from the zero external field and zero momentum value of the FWHM  $w$  of the composite peak (not shown), since in that case we expect the splitting  $\delta$  to vanish [see Fig. 1(a)] and both peaks to lie perfectly on top of each other. This yields  $0.124 \pm 0.005 \text{ meV}$ .

This linewidth value can be compared to theory. ISB spin plasmons are expected to be immune against DP dissipation [18, 19]. Thus, owing to the very high mobility of the sample and the low working temperature, we expect the linewidth to be dominated by an intrinsic many-body effect, the spin Coulomb drag [22–24] (SCD). The SCD is caused by a friction of Coulomb origin between carriers

of opposite spin moving with different momentum. The ISB spin plasmon, where spin densities oscillate out-of-phase along the growth axis, provides an optimal scenario for the SCD [23]. A calculation of the corresponding linewidth within a local-density approximation yields an SCD linewidth of the order of a fraction of an meV, confirming that the dominant dissipation source is the SCD. As the latter is mainly due to the out-of-plane spin density oscillation, its  $\mathbf{q}$  dependence, for  $\mathbf{q}$  much smaller than the Fermi momentum, is weak and only to second order. Hence, we deconvolute all the  $w(B_{\text{ext}})$  curves using the experimentally determined 0.124 meV.

Figure 4(e) presents  $\delta(B_{\text{ext}})$  (symbols as in Fig. 3), obtained by deconvolution of the data of Fig. 3(a). Using Eq. (3), we can now evaluate the collective magnetic moment of the spin plasmons as  $\mu = \delta(B_{\text{ext}} = 0) / (2 |\mathbf{B}_{\text{SO}}^{\text{coll}}|)$ . This ratio is plotted in Fig. 4(b) (squares), for the various  $\varphi$  probed. It appears constant with  $\varphi$  within the experimental error. We deduce  $\mu = 28.8 \pm 0.7 \mu\text{eV T}^{-1} = (0.50 \pm 0.01) \mu_{\text{B}}$ .

The consistency of our interpretation of the data with the model of Eq. (3) is demonstrated in Fig. 4(e), which compares the experimental data points for  $\delta(B_{\text{ext}})$  with the relation  $\delta(B_{\text{ext}}) = 2\mu\sqrt{(B_{\text{ext}} + B_{\text{SO},\parallel}^{\text{coll}})^2 + B_{\text{SO},\perp}^{\text{coll}}{}^2}$  (lines), using the previously determined values of  $B_{\text{SO},\parallel}^{\text{coll}}$ ,  $B_{\text{SO},\perp}^{\text{coll}}$  and  $\mu$ . An excellent agreement is found, *without introducing any fitting parameters*.

We further validate our model by checking the  $q$ -dependence of  $\mu$ . We repeat the same experimental procedure for other values of  $q$ . Figure 4(c) presents the values of the minimum ( $\varphi = 90^\circ$ , open diamonds) and maximum ( $\varphi = 0^\circ$ , filled diamonds) modulus of  $\mathbf{B}_{\text{SO}}^{\text{coll}}$ . They appear proportional to  $q$ . Figure 4(d) shows the angular average of the magnetic moment  $\mu$  (squares). Interestingly,  $\mu$  turns out to be practically constant with  $q$ . This demonstrates that all of the SO effects are contained in  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$ , and that  $\mu$  is indeed the largest quantized projection of the intrinsic ISB spin plasmon magnetic moment  $\mathbf{M}$  onto the field direction. This point is confirmed by noting that  $\mu \approx 2 \frac{|g|\mu_{\text{B}}}{2}$ , that is,  $\mu$  is very close to twice the magnetic moment of a single electron (when considering the  $g$ -factor of bulk GaAs,  $g = -0.445$ ). This is consistent with the fact that an ISB spin plasmon involves transitions between two spin 1/2 states, i.e. excitations of spin magnitude 1. Our results thus show that the ISB plasmon maintains the spin magnitude of a single elementary excitation, while the many-body effects are absorbed in the collective magnetic field  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$ . Hence, the quantized projection of the plasmon magnetic moment can either be  $\pm\mu$  [ $m_{\pm}$  modes, see Fig. 1(c)] or 0 [ $m_{\parallel}$  mode, whose energy is unaltered by SO coupling].

To summarize the experimental part: we validated our fine structure model by demonstrating the internal consistency of measurements—with and without magnetic field—with Eqs. (2) and (3). We emphasize again that

this did not involve any adjustable parameters.

All elements are now in place to see how the collective SO magnetic field  $\mathbf{B}_{\text{SO}}^{\text{coll}}$  emerges from the  $\mathbf{k}$ -dependent single-particle magnetic fields  $\mathbf{B}_{\text{SO}}(\mathbf{k})$  given by Eq. (1). The ISB spin plasmon is a superposition of single-particle transitions from momentum  $\mathbf{k}$  in the first subband to  $\mathbf{k} + \mathbf{q}$  in the empty second subband. Thus, each electron-hole pair experiences a crystal magnetic field difference  $\Delta\mathbf{B}_{\text{SO}}(\mathbf{k}, \mathbf{q}) = \mathbf{B}_{\text{SO},2}(\mathbf{k} + \mathbf{q}) - \mathbf{B}_{\text{SO},1}(\mathbf{k})$ , given by

$$\frac{g\mu_{\text{B}}}{2} \Delta\mathbf{B}_{\text{SO}}(\mathbf{k}, \mathbf{q}) = \alpha_2 \begin{pmatrix} q_y \\ -q_x \end{pmatrix} + \beta_2 \begin{pmatrix} q_x \\ -q_y \end{pmatrix} + (\alpha_2 - \alpha_1) \begin{pmatrix} k_y \\ -k_x \end{pmatrix} + (\beta_2 - \beta_1) \begin{pmatrix} k_x \\ -k_y \end{pmatrix}, \quad (4)$$

where the subscript  $n = 1, 2$  refers to the subband index. With  $\alpha_1 = 3.5 \text{ meV \AA}$ ,  $\alpha_2 = 2.8 \text{ meV \AA}$ ,  $\beta_1 = 0.22 \text{ meV \AA}$ , and  $\beta_2 = 0.79 \text{ meV \AA}$ , we are able to reproduce the experimental data in Fig. 2(b)–(d) and Fig. 4(a)–(e) in a quantitatively accurate way (see lines), using a linear-response formalism based on time-dependent density-functional theory.

$\Delta\mathbf{B}_{\text{SO}}(\mathbf{k}, \mathbf{q})$  contains a  $\mathbf{k}$ -independent part, which is thus the same for all electron-hole pairs, and a  $\mathbf{k}$ -dependent part. The latter could have a disorganizing effect, causing DP dephasing. This is indeed what occurs for single-particle spin dynamics [5, 8–11]. But here, the  $\mathbf{k}$ -dependence turns out to be exactly canceled by an additional, dynamical Coulombic contribution [19], explaining how a uniform  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$  can emerge.

In a simple scenario, one could expect  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q})$  to be aligned with the  $\mathbf{k}$ -independent part of  $\Delta\mathbf{B}_{\text{SO}}(\mathbf{k}, \mathbf{q})$ , with a slightly enhanced magnitude. But what is found is that  $\mathbf{B}_{\text{SO}}^{\text{coll}}(\mathbf{q}) = \frac{2 \times 5.25}{g\mu_{\text{B}}} (\bar{\alpha}q_y + \bar{\beta}q_x, -\bar{\alpha}q_x - \bar{\beta}q_y)$  (within 3%), with  $\bar{\alpha} = (\alpha_1 + \alpha_2)/2$  and  $\bar{\beta} = (\beta_1 + \beta_2)/2$ . That is, many-body effects tilt the  $\mathbf{k}$ -independent part of  $\Delta\mathbf{B}_{\text{SO}}(\mathbf{k}, \mathbf{q})$ , align it with the average single-particle SO field difference, and amplify it by about a factor five.

Such a magnification effect due to dynamical many-body interactions is quite remarkable. At first glance, one would expect Coulomb-induced enhancements to be roughly of order  $r_s$  (Wigner-Seitz radius), which is  $\simeq 1.3$  for the studied sample. On the other hand, recent experiments [25, 26] suggest that the interplay of Coulomb and SO interactions could manifest in a mutual boost, leading to significant enhancement of electronic spin splittings especially in low-dimensional systems [27].

In conclusion, we have shown that many-body effects can produce significant departure from the single-particle picture of SO effects in periodic solids. In an ISB spin plasmon, despite the spread of electronic velocities, a well-organized spin dynamics emerges at the collective level. The electrons coherently precess about a giant SO field, which gives rise to a fine structure of the spin-plasmon spectrum. This effect, which might also play a role in other helical liquids [27, 28], reveals novel op-

portunities for magnetization control with collective SO fields.

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\* Corresponding author: [florent.baboux@insp.upmc.fr](mailto:florent.baboux@insp.upmc.fr)

† Present address: Centro Atómico Bariloche, Bariloche, Argentina.

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# Supplemental Material for: Giant Collective Spin-Orbit Field in a Quantum Well: Fine Structure of Spin Plasmons

F. Baboux,<sup>1,\*</sup> F. Perez,<sup>1</sup> C. A. Ullrich,<sup>2</sup> I. D'Amico,<sup>3</sup> J. Gómez,<sup>1,†</sup> and M. Bernard<sup>1</sup>

<sup>1</sup>*Institut des Nanosciences de Paris, CNRS/Université Paris VI, Paris 75005, France*

<sup>2</sup>*Department of Physics and Astronomy, University of Missouri, Columbia, Missouri 65211, USA*

<sup>3</sup>*Department of Physics, University of York, York YO10 5DD, United Kingdom*

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## EXPERIMENT

The studied sample is an asymmetrically modulation-doped GaAs/Al<sub>0.3</sub>Ga<sub>0.7</sub>As quantum well, grown along the [001] direction by molecular beam epitaxy. The 200 Å-thick well was doped with two Si delta layers, separated from the well by spacer thicknesses of 350 and 400 Å respectively. The electron density is  $2.3 \times 10^{11} \text{ cm}^{-2}$ , and the mobility  $2 \times 10^7 \text{ cm}^2 \text{ V}^{-1} \text{ s}^{-1}$  at the working temperature  $T \simeq 2 \text{ K}$  (superfluid Helium), as determined from Hall measurements.

The inelastic light scattering setup, depicted in Fig. 2(a) of the Letter, is in the backscattering geometry. The average angle  $\theta$  of the incoming and backscattered light with respect to the normal direction can be changed to transfer a Raman momentum  $\mathbf{q}$  of amplitude  $q = \frac{4\pi}{\gamma} \cos \frac{\beta}{2} \sin \theta$ , with  $\gamma \simeq 3^\circ$  and  $\lambda \simeq 770 \text{ nm}$  the exciting wavelength.

Measurements are performed in an pumped-Helium optical cryostat embedding a superconducting coil. The scattered light is dispersed by a Dilor XY triple Raman spectrometer in additive mode.

To deconvolute the linewidth  $w$  of the composite Raman peak observed (Fig. 2(a) of the Letter), we fit the sum of two identical Lorentzian peaks of linewidth  $\Gamma$  separated by  $\delta$ , with a single Lorentzian, yielding a linewidth  $w$ . We then numerically invert the  $w(\delta)$  plot to obtain the  $\delta(w)$  function.

## CALCULATION OF INTERSUBBAND SPIN PLASMON DISPERSIONS

The formal framework for calculating intersubband spin plasmon dispersions in quantum wells, in the presence of spin-orbit coupling, is based on time-dependent density-functional theory within the effective-mass approximation, as described in detail in Refs. [1, 2]. We have here extended this formalism to account for external in-plane magnetic fields. The calculation proceeds in two steps. First, the two-component spinor subband

envelope functions  $\psi_j(\mathbf{k}, z)$  are calculated via

$$\sum_{\beta=\uparrow,\downarrow} \left[ \hat{h} \delta_{\alpha\beta} + v_{\alpha\beta}^{\text{xc}} + \frac{g\mu_B}{2} [\vec{\sigma} \cdot (\mathbf{B}_{\text{ext}} + \mathbf{B}_{\text{SO}}(\mathbf{k}))]_{\alpha\beta} \right] \psi_{j\beta} = \varepsilon_j \psi_{j\alpha}, \quad (1)$$

where  $\hat{h}$  is the single-particle kinetic energy operator plus quantum well confining potential and Hartree potential,  $v_{\alpha\beta}^{\text{xc}}$  is the spin-dependent exchange-correlation (xc) potential in local-density approximation, and  $\vec{\sigma}$  is the vector of Pauli matrices. The parameters of our modulation-doped GaAs quantum well are taken as input.

The second step is to calculate the intersubband spin plasmon dispersions using linear-response theory. Formally, the plasmon excitation energies are obtained from the poles of the response function

$$\Pi = \frac{\Pi_0}{1 - \Pi_0 F^{\text{Hxc}}}, \quad (2)$$

where  $\Pi_0$  is the noninteracting response function, built from the solutions of Eq. (1), which determines the single-particle excitation spectrum.  $\Pi_0$  and  $\Pi$  are represented as  $4 \times 4$  matrices, describing the charge-density response and three components of the magnetization response. The formation of collective plasmon excitations is a consequence of dynamical many-body effects (Hartree and xc), accounted for via  $F^{\text{Hxc}}$ . The spin plasmons arise solely from the xc part of  $F^{\text{Hxc}}$ , which is treated in an adiabatic local-density approximation [1, 2].

Based on the given quantum well parameters, the Rashba and Dresselhaus constants for the  $n$ th subband,  $\alpha_n$  and  $\beta_n$ , can be calculated using  $\mathbf{k} \cdot \mathbf{p}$  theory[3, 4]. For the Dresselhaus parameter for bulk GaAs we instead use a recent result from an ab-initio GW calculation [5],  $\gamma_{\text{GaAs}}^{\text{GW}} = 6.4 \text{ eV \AA}^3$ , which is about 4 times smaller than the  $\mathbf{k} \cdot \mathbf{p}$  value. We thus obtain  $\alpha_1^{\text{kp}} = 0.84 \text{ meV \AA}$ ,  $\alpha_2^{\text{kp}} = 0.66 \text{ meV \AA}$  and  $\beta_1^{\text{GW}} = 0.93 \text{ meV \AA}$ ,  $\beta_2^{\text{GW}} = 3.3 \text{ meV \AA}$ . However, to reproduce the experimental data in Figs. 2(b)–(d) and 4 of the Letter, we found it necessary to rescale the values of the Rashba and Dresselhaus constants by 4.2 and 0.24, respectively, i.e., to use  $\alpha_{1,2} = 4.2\alpha_{1,2}^{\text{kp}}$  and  $\beta_{1,2} = 0.24\beta_{1,2}^{\text{GW}}$ . We emphasize that these are the only fitting parameters of our theory. Note that  $\sqrt{\alpha^2 + \beta^2} \simeq 1.4 \times \sqrt{\alpha^{\text{kp}2} + \beta^{\text{GW}2}}$ . Thus, the enhancement factor from single-particle to collective spin-orbit magnetic field, which is found of 5.25

with the fitted values for  $\alpha$  and  $\beta$ , would be close to 7.3 with the predicted ones.

### CALCULATION OF THE SPIN COULOMB DRAG LINEWIDTH

The spin Coulomb drag linewidth  $\Gamma_{\text{SCD}}$  is calculated within the framework of time-dependent spin-current density-functional theory, with dissipative effects (SCD and viscosity) included from first principles, see Ref. [6]. The SCD enters the formulation as one of the contributions to the exchange-correlation kernel. This contribution is calculated within a three-dimensional local-density approximation, which accounts for inhomogeneity in the growth direction.

Improving over Ref. [6], we now accurately model the first-subband envelope function using a density-functional scheme, with the parameters of the experimental quantum well (doping density, quantum well width, doping-layer position) as input. As the SCD linewidth may be very sensitive to the envelope function shape, the uncertainty of these input parameters is important. For the sample considered, reasonable uncertainties over the experimental parameters produce a SCD linewidth variation of about 15%. We have also derived the finite- $q$  correction to  $\Gamma_{\text{SCD}}$ , which is only second order in  $q$ . In contrast with Ref. [6], we have included in the current calculations the dissipative effects stemming from the viscosity term of the exchange-correlation kernel. In principle, this contribution might be substantial as the quantum well is strongly asymmetric, but it turns out to be one order of magnitude smaller than the SCD contribution. This confirms the latter to be the largest source of intrinsic dissipation for an ISB spin plasmon.

For the studied sample we find  $\Gamma_{\text{SCD}} \simeq 0.4$  meV, which

overestimates the linewidth. We find a quantitatively similar overestimate when calculating the spin plasmon linewidth of the experimental results in Ref. [7]. In the case of charge plasmons in quantum wells it was argued [8] that the three-dimensional local-density approximation does not account properly for the subband quantization, which acts as a bottleneck for energy-momentum dissipation from the collective to the in-plane degrees of freedom. Hence, it is not surprising that our theoretical estimate for  $\Gamma_{\text{SCD}}$  lies above the experimental linewidth while still providing the correct order of magnitude.

There are no adjustable parameters in the theory.

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- \* Corresponding author: florent.baboux@insp.upmc.fr  
 † Present address: Centro Atómico Bariloche, Bariloche, Argentina.
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