Spin susceptibility enhancement in a two-dimensional hole gas

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The critical spin splitting needed to fully polarize a two-dimensional heavy-hole gas in a (Cd,Mn)Te quantum well thanks to the giant Zeeman effect is measured through photoluminescence and transmission spectroscopy. While the splitting between two circularly polarized peaks in photoluminescence remains equal to the bare spin splitting, the critical spin splitting which governs the polarization is enhanced, in good agreement with the prediction of a model of hole-hole interactions previously applied to conduction electrons. Consequences for carrier-induced ferromagnetism are discussed.

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I. INTRODUCTION

A considerable amount of work has been devoted during the past 30 years to the study of two-dimensional systems of charge carriers, with the mainstream focused on electrons in semiconductor quantum wells and heterojunctions. The importance of carrier-carrier interactions in a two-dimensional system was recognized early[1] and intensively studied since then in two-dimensional electron gases (2DEG). In spite of that, this is still a very active area of research, and a complete understanding has not been achieved yet.

One of the consequences of carrier-carrier interactions is an enhancement of the spin susceptibility.[2–4] The spin polarization of a noninteracting 2DEG with a parabolic in-plane dispersion involves two material parameters: The effective mass describing the in-plane motion, m^*, which enters the density of states $m^*/\pi\hbar^2$, and the Landé factor g, which enters the spin splitting $g\mu_B\mu_0 H$ due to an applied field H. At low temperature (i.e., k_BT much smaller than the Fermi energy E_F) the resulting Pauli susceptibility χ_0 is a constant, which does not depend on the applied field or on the spin polarization. In the case of an interacting 2DEG, the effect of interactions on the spin susceptibility has been characterized experimentally by transport measurements.[5–8] Experimental data and theoretical studies all conclude that the resulting spin susceptibility χ is enhanced, but there is no agreement on a general description of this enhancement for all values of the carrier density and spin polarization.

The enhancement of the spin susceptibility should have a direct effect on carrier-induced ferromagnetism in semiconductors.[9,10] This is a specific property of diluted magnetic semiconductors which has recently boosted the interest in these materials, especially because applications are contemplated in the field of spintronics. Localized spins, due to magnetic impurities substituting a constituent atom of a normal semiconductor, experience a ferromagnetic interaction induced by free carriers. More precisely, the two systems (localized spins on one hand and free carriers on the other hand) are coupled by a Kondo-like Hamiltonian which has been deduced from experimental studies of the giant magneto-optical properties in II-VI semiconductors.[11,12] In the simplest model, disorder is neglected and the effect of each system on the other is treated in a double mean-field approximation. Then, the critical temperature is calculated to be proportional to the Pauli susceptibility of the carriers.[13,14] Carrier-induced ferromagnetism has been demonstrated in several three-dimensional (3D) systems (thick epitaxial layers).[13,15,16] Another very interesting configuration is that of carrier-induced ferromagnetism in 2D systems, that has been demonstrated in p-type modulation doped $Cd_{1-x}Mn_xTe$ quantum wells[17,18] and more recently in $Ga_{1-x}Mn_xAs$ heterojunctions.[19] As the spin-carrier coupling is larger in the valence band than in the conduction band, one has to deal with a 2D hole gas (2DHG). An enhancement of the susceptibility of a 2DHG in a $Ga_{1-x}Mn_xAs$ quantum well (QW), by a factor of about 2, was found theoretically using models which do not take into account the specific character of the valence band.[14,20] The same value of this factor, that leads to an enhanced critical temperature,[14,17,21] has been used in the case of $Cd_{1-x}Mn_xTe$ QWs.[9,17]

The case of a 2DHG is actually more difficult to handle than the 2DEG, because of the orbital degeneracy of the top of the valence band that has a p-like symmetry, the resulting anisotropy of the effective mass of holes, and the effect of spin-orbit coupling: The spin polarization has to be defined precisely and cannot be assimilated to a mere population distribution on spin states.[13,22] To give just a few examples, this has consequences on various properties such as transport properties and the presence of a metal-insulator transition,[23] Landau levels and quantum Hall effect,[24,25] and elementary excitations of the 2DHG.[26] This makes the 2DHG an interesting system in itself. Carrier-carrier interactions have to be considered in addition to this already intricate situation.

An experimental determination of the enhancement of the spin susceptibility by carrier-carrier interactions in a 2DHG is thus of importance, and particularly in the frame of carrier-induced ferromagnetism. It should be carried on preferably in systems, where carrier-induced ferromagnetism has been reported, and where the material parameters are known. A spectroscopic determination is particularly attractive, since it is direct and explicit. Kohn's theorem,[27] or its spin version the Larmor theorem,[28] states that optical transitions are not sensitive to carrier-carrier interactions, when the carriers have the same effective mass. So, optical transitions at $k=0$ cannot discriminate between single particle or collective excitations. While it is in principle established for intraband transitions, we check in this work if this is valid even for photoluminescence transitions related to charged excitons. Two-photon spectroscopy, such as Raman spectroscopy, escapes this limitation. It has been applied to CdTe and $Cd_{1-x}Mn_xTe$ based QWs containing a 2D electron gas,[29] and earlier to GaAs QWs with a 2DHG.[30] No results exist up to now in the case of a $Cd_{1-x}Mn_xTe$ QW with a 2DHG, which is the model system for 2D carrier-induced ferromagnetism.

In this work we use photoluminescence (PL) spectroscopy to measure the enhancement of the spin susceptibility by hole-hole interaction in $Cd_{(1-x)}Mn_xTe$ QWs. We first measure the giant Zeeman splitting $Z_h(H)$ at low carrier density. Then, we confirm that the same value of the bare Zeeman splitting $Z_h(H)$ determines the splitting between PL lines of opposite circular polarization in the high density regime, in agreement with the Larmor theorem. However, the PL spectra exhibit a specific feature that, in addition to the comparison of the PL and transmission peak positions, allows us to identify the value of the Zeeman splitting for which the gas is fully polarized. We show that this critical splitting $Z_h(H_c)$ is definitely smaller than the splitting $Z_h^*(H_c)$ which would be needed in the noninteracting 2DHG. In other words, we have an effective splitting $Z_h^*(H_c)$, which is enhanced with respect to the actual applied splitting $Z_h(H_c)$. We then compare our experimental value of the enhanced effective splitting to the prediction of a recent model,[31] which calculates the spin susceptibility of the interacting 2DHG, $\chi(p,\zeta)$, as a function of the density p and polarization degree ζ. The enhancement of the critical effective splitting and that of the spin susceptibility are linked by $Z_h(H_c)/Z_h^*(H_c)=\int_0^1(\chi_0/\chi)d\zeta$ (see Sec. IV). Finally, we discuss the influence of this enhancement on carrier-induced ferromagnetism in QWs.[17]

II. SAMPLES AND EXPERIMENTAL DETAILS

Samples were $Cd_{1-x}Mn_xTe$ QWs with a low Mn content ($x<0.01$), grown by molecular beam epitaxy on a $Cd_{1-z}Zn_zTe$ substrate (with $z\approx 0.12$). A 2DHG is generated by p-type modulation doping with nitrogen impurities inserted in the $Cd_{1-y-z}Zn_yMg_zTe$ barriers (typically $y\approx 0.08$ and $z\approx 0.25$),[32] or induced by electron trapping on acceptor states at the surface of the sample.[33] The splitting between the light and heavy holes due to confinement in the 8 nm-wide QW, is made even larger by the biaxial strain resulting from the coherent growth. The calculation of the hole splitting uses known values of the deformation potential[34] and of the Luttinger parameters in CdTe.[35] It has been validated in samples specially designed to feature the light-hole/heavy-hole crossing,[36] and in all the present samples it predicts a splitting larger than 30 meV. Since the Fermi energies of the 2DHG which are investigated here are of a few meV, only the heavy-hole subband is populated. Similar QWs, but with a larger Mn content ($x\leqslant 0.03$), exhibit carrier-induced ferromagnetism.[17,18]

Details on the optical set up, and on the determination of the Mn content and carrier density, are given in Refs. 37 and 38. We only give a brief summary here. It was possible to perform experiments at different carrier gas densities in the same sample, by illuminating the sample with photons of energy larger than the band gap of the barrier. The mechanism is an intrinsic one: Electrons created in the barrier drift to the QW due to the built-in electric field, while holes slowly tunnel through the triangular barrier due to the same electric field.[37] The carrier density p was determined using a calibration[38] that relies on an extensive use of several methods (such as Hall effect, filling factors of Landau levels at high field) independent of material parameters, and Moss-Burstein shift which depends on the electron and hole effective mass. Actually, most of the samples used in the present study have been used also in Ref. 38, thus ensuring a good determination of the hole density.

III. OBSERVATION OF THE BARE ZEEMAN SPLITTING IN SPECTROSCOPY

Spectra related to interband transitions in such QWs have been described in detail in Refs. 37 and 38. In the region of moderately giant spin splitting, which is relevant in the present study, we have shown[38] that the PL and transmission (or reflectivity, or PL excitation) lines are related to the annihilation/creation of a charged exciton X^+.

At vanishingly low carrier density, the charged exciton is a three-particle complex, containing two holes (in a singlet configuration), and one electron. One hole preexists in the QW, the electron and the second hole are photocreated. If the preexisting holes are spin-polarized (say $+\frac{3}{2}$ holes in a (Cd,Mn)Te QW with a magnetic field applied in the Faraday configuration), the charged exciton is created by the absorption of a photon with a σ^- helicity (the photocreated hole is $-\frac{3}{2}$ and the electron $+\frac{1}{2}$). In PL, such a charged exciton emits a photon of the same helicity. However, if the electron flips the spin from $+\frac{1}{2}$ to $-\frac{1}{2}$ (which, in absence of electron gas, costs an energy equal to the giant Zeeman splitting Z_e in the conduction band), the charged exciton recombination produces a σ^+ photon and leaves a $-\frac{3}{2}$ hole. In the case of a partially polarized hole gas of vanishingly low density, each transition (σ^- and σ^+) is observed both in PL and in transmission, at the same energy (but for a small Stokes shift). The initial state of the PL transition is the charged excitons with two holes in a singlet state and an electron with spin up or down: The splitting of these initial states for the σ^- and σ^+ transitions is the giant Zeeman splitting of the electron, $Z_e(H)$. The final state is a hole with spin up or down: The splitting between the final states is that of the hole, $Z_h(H)$. Thus the splitting between the two transitions is the sum $Z_X(H)=Z_e(H)+Z_h(B)$. The experimental splitting is well reproduced by a Brillouin function, and it is known[12] that, in (Cd,Mn)Te, $Z_h(H)/Z_e(H)=4.0$. A symmetrical shift of the σ^- and σ^+ lines is observed as long as the giant Zeeman splitting remains small with respect to the QW effective depth (energy

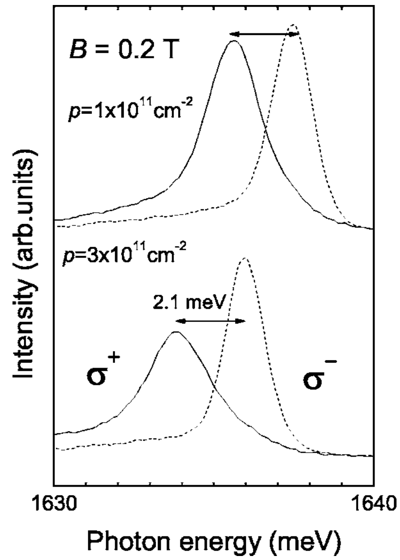

FIG. 1. PL spectra of a $Cd_{0.996}Mn_{0.004}Te$, 8 nm wide QW, in σ^+ polarization (solid line) and σ^- polarization (dashed line), for two values of the carrier density, $p=1\times10^{11}$ cm^{-2} (top) and $p=3\times10^{11}$ cm^{-2} (bottom).

gap from the confined levels to the barrier), which is the case in the present study.

The description of the transitions associated to the charged exciton have to be reconsidered at larger values of the carrier density, such that the Fermi energy of the hole gas is larger than the thermal energy and characteristic energies governing the localization of the carriers.[37,38] Figure 1 shows PL spectra in σ^+ and σ^- polarization, for two different values of the carrier density in the same QW, respectively 1×10^{11} cm^{-2} and 3×10^{11} cm^{-2}. Surprisingly, the main changes are a broadening of the line and a redshift, without significant change of the σ^+/σ^- splitting. We first focus on the splitting and we will come back to the position of the lines [Figs. 2(a) and 2(b)] in the next section.

Figure 2(c) displays the σ^+/σ^- splitting as a function of the applied field, for several values of the carrier density. It exhibits only a small change, if any, when the carrier density increases from values where the spectra are characterized by sharp neutral exciton and charged exciton lines, to the maximum value $p=3\times10^{11}$ cm^{-2} in the present sample. It can be noticed in Fig. 1 that the small changes are indeed much smaller than the line width, which increases when the carrier density increases. The broadening in doped QWs, which exhibits a low energy tail, is probably the result of the creation of other excitations in the 2DHG during the PL transition, as calculated by Esser[39] for the absorption line of the charged exciton transition in a quantum wire.

We conclude that whatever the hole density, the value of the σ^-/σ^+ splitting remains equal to the sum of the bare Zeeman splitting, $Z_X(H)$, previously determined for vanishing density. This experimental fact is a key argument to support the result of this work.

We will show below that the polarization of the hole gas is governed by an effective splitting $Z_h^*(H)$ which is enhanced with respect to $Z_h(H)$, and can be considered as the energy needed to flip the spin of a single hole. In the final state of the σ^- PL transition, the hole gas has the same polarization as the ground state, while in the final state of the σ^+ transition one hole has been transferred from the majority spin subband to the minority subband. One could then expect the σ^-/σ^+ splitting to be $Z_e(H)+Z_h^*(H)$: This is not supported by our experimental result. A more convenient interpretation is that the recombination process involved in the σ^+ PL transition implies the emission of a collective spin-flip excitation of the hole gas with vanishing wave vector. It is indeed well known, at least for electrons that, according to Larmor's theorem, the energy of the zone center collective spin-flip excitation is exactly the bare Zeeman energy [say, for holes, $Z_h(H)$] whatever the carrier density.[40] Then the expected splitting is $Z_e(H)+Z_h(H)$, as actually observed.

IV. ENHANCEMENT OF THE SPIN SUSCEPTIBILITY

Figures 2(a) and 2(b) display the position of the PL and absorption lines as a function of the applied field, for the two

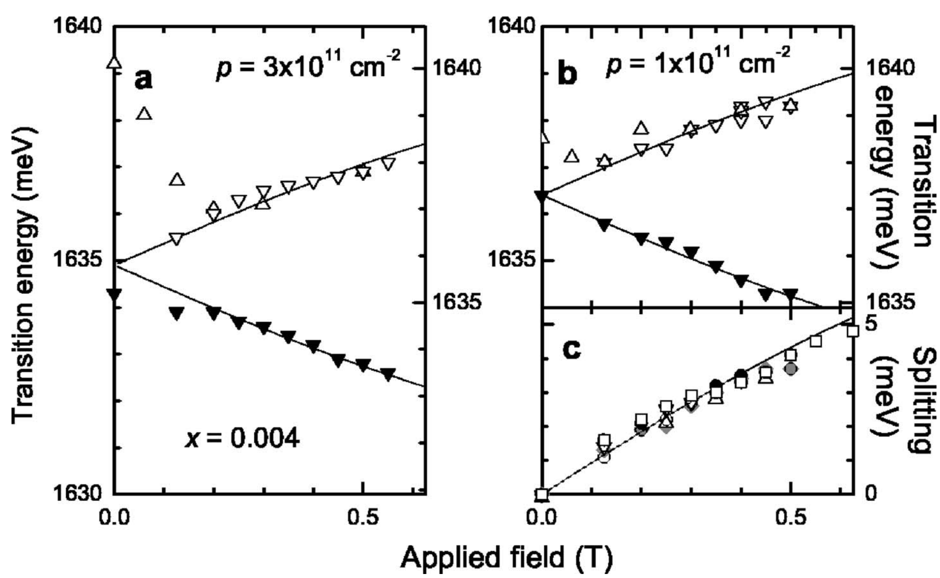

FIG. 2. Transition energies in PL (down triangles, left scale) and transmission (up triangles, right scale) in σ^+ polarization (full symbols) and in σ^- polarization (open symbols) for two values of the carrier density, (a) $p=3\times10^{11}$ cm^{-2}, and (b) $p=1\times10^{11}$ cm^{-2}; (c) Splitting between the σ^+ and σ^- PL lines, as a function of the applied field, for different values of the carrier density from $p=0.7\times10^{11}$ cm^{-2} (open triangles) to 3×10^{11} cm^{-2} (full circles). The solid lines show the Brillouin function describing the effect of the giant Zeeman effect in $Cd_{0.996}Mn_{0.004}Te$.

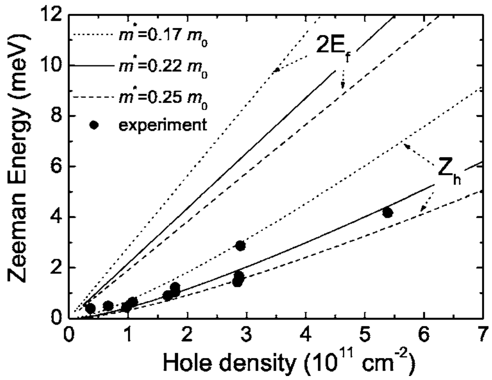

FIG. 3. Valence band spin splitting achieving complete polarization of the 2DHG (symbols: experimental data; lower set of lines: calculated), and Fermi energy (upper set of lines), as a function of the carrier density, assuming three values of the transverse mass of the heavy holes, $m^*=0.17m_0$ (diagonal approximation, dotted lines), $m^*=0.25m_0$ (axial approximation, dashed lines) and $m^*=0.22m_0$ (best fit, solid lines).

values of the hole density used in Fig. 1. From Fig. 2(a), one clearly distinguishes two field ranges for the highest carrier density achieved in this sample. Above 0.25 T, the σ^- lines coincide in PL and transmission (except for a small, constant Stokes shift, which has been taken into account by a shift of the two vertical scales in Fig. 2), and the σ^+ PL line exhibits a symmetrical shift. The shifts are well reproduced by the Brillouin function. The coincidence of the two lines in σ^- polarization strongly suggests that they involve the same states; in particular, the final state in σ^- PL is the same as the initial state in σ^- transmission, i.e., it is the ground state. In this field range, the spin splitting is large enough that the hole gas be fully polarized: all holes are $+\frac{3}{2}$. As discussed in the previous section, the σ^+ PL transition involves the recombination of the $-\frac{1}{2}$ electron with a $+\frac{3}{2}$ hole, thus leaving a collective spin-flip excitation in the hole gas.

Below 0.2 T, all line shifts deviate from the Brillouin function, and a kind of a Moss-Burstein shift appears between the σ^- lines in PL and in transmission. We have already shown that, in spite of the role played by charged excitons in the optical transitions for this range of spin splitting, this shift is proportional to the density of carriers with the relevant spin.[38] We have argued that, for both polarizations, the transitions involve excitations of the system in addition to the creation or annihilation of the charged exciton. Here we simply use the onset of the Moss-Burstein shift in σ^- polarization as a measure of the smallest value of the applied field, H_c, needed to fully polarize the 2DHG.

The same effects are observed at lower hole density [Fig. 2(b)] with a much smaller value of the characteristic field where the Moss-Burstein shift appears. If the carrier density is further reduced, spectra are characterized by sharp lines related to the charged and neutral exciton that follow the Brillouin function.[37] We do not consider this density range in the present work.

The valence band bare spin splitting at full polarization, $Z_h(H_c)$, is shown in Fig. 3 as a function of the carrier density. In the absence of carrier-carrier interactions, we would expect $Z_h(H_c)=2E_F$ where E_F is the Fermi energy of the unpolarized hole gas (upper curves).

However, the determination of E_F is not obvious: Even if the carrier density has been thoroughly determined, we have no direct determination of the Fermi energy. In particular, the Moss-Burstein shift involves a contribution from both bands, and the contribution from the conduction band is larger than the contribution from the valence band. In Fig. 3, three values of the transverse mass of the heavy holes have been used to calculate the Fermi energy. In the so-called diagonal approximation, $m^*=(\gamma_1+\gamma_2)^{-1}$, where γ_1 and γ_2 are the relevant Luttinger parameters deduced from cyclotron resonance in CdTe,[35] we obtain $m^*=0.17m_0$. The Moss-Burstein shift measured in Ref. 38 is too small to support such a small value. A more reasonable choice is to use the axial approximation[42] which leads to $m^*=0.25m_0$ for the width and strain values of these samples. Finally, the best fit to the data was obtained for $m^*=0.22m_0$, which is very close to the previous value.

As a sign of the enhanced susceptibility, we observe a value of $Z_h(H_c)$ which is much smaller than $2E_F$, regardless of the mass we used to calculate E_F.

Indeed, exchange and correlations Coulomb interactions are predicted to enhance the static spin susceptibility.[2–4] The density dependence of this enhancement has already been evidenced by transport measurements for conduction electrons confined in a GaAs/GaAlAs heterostructure[5,7,8] and for low spin-polarization rate. The values measured for spin-polarization rates below 20% have been found to be in good agreement with a recent calculation valid for vanishingly small polarization rate.[4] Here, we deal with holes and high spin-polarization rates (100%). As we expect a dependence of the static spin susceptibility on the polarization rate, we have derived χ for arbitrary polarization rate and we have considered the Zeeman energy enhancement factor Z_h^*/Z_h which can be extracted from our measurements. Calculations have been done for electrons and translated to the heavy-hole gas as justified below.

It is well known that a spin-polarized system presents a strong electron spin resonance (ESR)[43,44] when excited by the magnetic field of a microwave oscillating in the plane perpendicular to the quantizing static magnetic field. This resonance is due to the absorption of the microwave by the collective spin-flip wave of the gas and does occur when the energy of the microwave photon matches the Zeeman splitting imposed on the gas, regardless of the electron density. This fact is a consequence of the Larmor's theorem which states that the collective motion of electron spin is sensitive to the external field only,[28] or identically speaking, that the collective spin-flip wave energy at vanishing wave vector is equal to the bare Zeeman splitting Z. The frequency of the spin-flip wave appears as a pole of the dynamical spin-flip susceptibility which in the local spin-density approximation is given by[40,41]

$$\chi_+(\vec{q},\omega)=\frac{\Pi_{\uparrow\downarrow}(\vec{q},\omega)}{1-\dfrac{2}{p^2}\dfrac{1}{\zeta}\dfrac{\partial E_{xc}}{\partial\zeta}\Pi_{\uparrow\downarrow}(\vec{q},\omega)}, \quad (1)$$

where p is the hole density, ζ the spin polarization degree, E_{xc} the exchange-correlation part of the ground state energy and $\Pi_{\uparrow\downarrow}(\vec{q},\omega)$ is the spin-flip noninteracting polarization function[41]

$$\Pi_{\uparrow\downarrow}(\vec{q},\omega)=\frac{1}{S}\sum_{\vec{k}}\frac{p_{\vec{k}+\vec{q},\uparrow}-p_{\vec{k},\downarrow}}{\varepsilon_{\vec{k}+\vec{q},\uparrow}-\varepsilon_{\vec{k},\downarrow}-\hbar\omega}. \quad (2)$$

Here, S is the sample surface, $p_{\vec{k},\sigma}$ is the occupation number of the single heavy-hole energy $\varepsilon_{\vec{k},\sigma}$, which is the usual Fermi distribution at 0 Kelvin for a spin σ

$$\varepsilon_{\vec{k},\sigma}=\frac{\hbar^2k^2}{2m^*}+sgn(\sigma)\frac{Z_h^*}{2}, \quad (3)$$

where $\mathrm{sgn}(\sigma)$ equals +1 for spin up and −1 for spin down holes. For vanishing wave vector q, the polarization function reduces to

$$\Pi_{\uparrow\downarrow}(\vec{q},\omega)=\frac{m^*}{\pi\hbar^2}\frac{Z_h^*}{Z_h^*-\hbar\omega}. \quad (4)$$

The pole of $\chi_+(\vec{q},\omega)$ gives the spin-flip wave frequency and leads to the relation between Z_h and Z_h^* which is expressed as

$$\frac{Z_h}{Z_h^*}(p,\zeta)=1+\frac{m^*}{\pi\hbar^2p^2}\frac{1}{\zeta}\frac{\partial E_{xc}}{\partial\zeta}. \quad (5)$$

Analytic expressions of E_{xc} are not available, but Attacalite *et al.*[4] deduced, from quantum Monte Carlo calculations including Coulomb exchange and Coulomb correlations, parametrized expressions valid for any ζ. Equation (1), initially built for electrons,[45] has been compared with transport data of Vakili *et al.*,[5] and with Raman spectroscopy measurements;[31] a very good agreement has been obtained.

We have translated this formalism to heavy-holes and one may ask why a model designed for conduction electrons should also apply to holes. In these samples, this is justified by the light- and heavy-hole energy separation and by the fact that they do not mix when a magnetic field is applied perpendicular to the 2DHG plane. The magnetic field, applied along z, splits the heavy-hole doublet by a few meV, thanks to the giant Zeeman effect. So, the splitting of the heavy-hole doublet remains small with respect to the light-hole/heavy-hole splitting. One important consequence of this configuration is the quenching of the in-plane spin components. In a simplified description which uses the language of spherical symmetry, the heavy-holes states at $k=0$ are made of an orbital state of p-like symmetry, and a $-\frac{1}{2}$ spin state, both quantized along the normal z to the QW, i.e. $|\pm\frac{3}{2}\rangle=|\pm1\rangle|\pm\frac{1}{2}\rangle$. At the lowest order, the diagonal elements, such as those of S_z, stay unchanged, but off-diagonal elements, such as those of S_x or S_y, are fully quenched by the product of orbital states. This quenching of the transverse spin components implies that the off-diagonal terms of the ground state density matrix $n_{\uparrow\downarrow}$ and $n_{\downarrow\uparrow}$ are set to zero. We stress that this assumption realizes a full match between such a heavy-hole gas and the homogeneous electron gas. As a consequence, the ground state energy E_{xc} given by the quantum Monte Carlo calculations, the spin-flip susceptibility written as (1) and the Larmor's theorem apply to this heavy-hole gas. This suggests that the heavy-hole spin susceptibility enhancement is the same type of the electron spin susceptibility enhancement observed in Refs. 5, 7, 8, and 47. Of course, this would not hold for a hole gas of higher density or smaller light-hole/heavy-hole splitting, as in the QWs considererd by Winkler *et al.*[46]

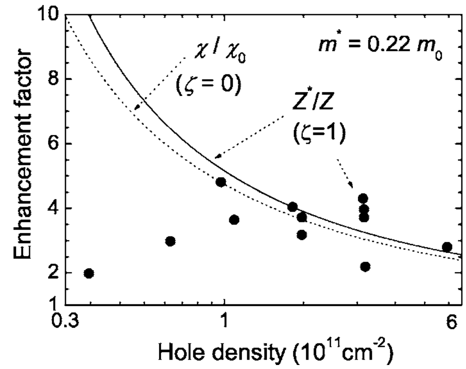

FIG. 4. Enhancement factor of the spin susceptibility and spin splitting, as a function of the carrier density, using $m^*=0.22m_0$; the lines give the result of the calculation at complete polarization (solid line) and vanishing polarization (dashed line); symbols are experimental data.

The experimental and calculated enhancement of the Zeeman splitting at full polarization, $Z_h^*(p,\zeta=1)/Z_h$, and the spin susceptibility enhancement at vanishing polarization, $\chi^*(p,\zeta=0)/\chi_0$, are compared in Fig. 3. The first quantity is measured in this work and the second is needed in the mean-field model of carrier-induced ferromagnetism. Both quantities are related as explained below.

The spin susceptibility of the interacting gas is the derivative of the spin density $(p\zeta)$ with respect to the applied field H, at constant carrier density: $\chi=d(p\zeta)/dH$. Note that here the driving force of the polarization is the spin splitting Z_h due to the giant Zeeman effect, instead of the applied field in the case of normal Zeeman effect. The spin susceptibility can be written as

$$\chi(p,\zeta)=p\frac{d\zeta}{dZ_h}\frac{dZ_h}{dH}=\frac{p}{2E_{\mathrm{F}}}\frac{dZ_h^*}{dZ_h}\frac{dZ_h}{dH}, \quad (6)$$

and the noninteracting spin susceptibility is defined as

$$\chi_0=\frac{p}{2E_{\mathrm{F}}}\frac{dZ_h}{dH}. \quad (7)$$

Thus, we obtain for any polarization degree

$$\frac{\chi(p,\zeta)}{\chi_0}=\frac{dZ_h^*}{dZ_h}(p,\zeta). \quad (8)$$

As only the enhancement for the full polarized state is given by our experimental data, we better use the relation obtained by integrating Eq. (8) taken at full polarization,

$$\frac{Z_h}{Z_h^*(p,\zeta=1)}=\int_0^1\frac{\chi_0}{\chi(p,\zeta)}d\zeta. \quad (9)$$

As shown in Fig. 4, the two enhancement factors remain close to each other, to within 10% for hole densities between 0.6 and 8×10^{11} cm^{-2}, and less than 15% for the smallest densities experimentally available, which is in both cases below the experimental accuracy.

The enhancement factor is calculated to be slightly larger than 2 in the high 10^{11} cm^{-2} density range; this value was assumed in the case of samples exhibiting carrier-induced ferromagnetism[17] which however have a smaller carrier density,[38] about 3×10^{11} cm^{-2}. As shown in Fig. 4, the calculated enhancement follows a monotonous increase as the carrier density decreases, and a reasonably good agreement is found with the experimental data down to 1×10^{11} cm^{-2}. That means that the effect was underestimated in Refs. 17 and 21. The effect should be even larger at low density, when correlations dominate over exchange. The experimental data are significantly lower than the calculated curve, as already reported for GaAs QWs with electron densities in the 10^{10} cm^{-2} range.[7] This may be partly due to errors in the determination of Z_h, which suffers from larger uncertainty at low density and doesn't reach zero as it should (see Fig. 3). Actually, this is the range where changes in the spectroscopic properties are expected—such as appearance of the neutral exciton beside the charged exciton, a stronger influence of localization on the properties of the carrier gas, and also stronger effects of disorder in carrier-induced ferromagnetism.[21,48,49]

V. CONCLUSION

We have measured the spin splitting needed to fully polarize a 2DHG, using the giant Zeeman splitting in the diluted magnetic semiconductor (Cd,Mn)Te, and the onset of the Moss-Burstein shift. It is much larger than the value expected in the absence of hole-hole interaction. In the range of moderate hole density ($1-6\times10^{11}$ cm^{-2}), this enhancement well agrees with the enhancement of spin susceptibility by exchange and correlation interactions, calculated using a model previously developed for a 2DEG. Both the experimental and calculated values are slightly larger than the value previously assumed in the study of carrier-induced ferromagnetism in a (Cd,Mn)Te QW. In agreement with the Larmor's theorem, the splitting between two circularly polarized components of photoluminescence remains equal to the bare spin splitting, whatever the carrier density is.

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