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Magnetoacoustic Plasmons in a Bilayer Quasi-Two-Dimensional Spin-Polarized System

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We investigate the charge and spin response functions of a bilayer quasi-two-dimensional system, spin polarized by a constant magnetic field $\vec{B}$. Terms beyond the random-phase approximation, the exchange and correlation interactions, are introduced by using generalized spin-dependent local field factors, $G_\sigma^{x,c}(\vec{q},\omega)$. The self-consistent magnetic interaction among the electron spins determines the coupling of the charge and the longitudinal spin-density excitations, leading to coupled in-phase and out-of-phase electric and magnetic modes. We find that the lowest frequency belongs to an acoustic mode, that represents the out-of-phase oscillation of the longitudinal magnetizations in the two layers. This collective excitation is shown to become important in the case of materials with large gyromagnetic factors, such as dilute magnetic semiconductors.

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As a measure of the strength of the interaction in a many-particle environment, the departure of the collective modes from the single-particle excitation frequency has often been a conclusive probe of the microscopic properties of the system, easily bridging theory with experimental data obtained in spectroscopical studies.

Within the random phase approximation (RPA), the excitation frequencies of the collective modes emerge as the poles of the response functions to an electromagnetic perturbation and in certain cases they have been known for a long time. A single unpolarized 2D layer exhibits a single charge mode at plasma frequency $\omega_p \sim \sqrt{q}$ (Ref. 1) while a bilayer system was found to have two charge modes, associated with the in-phase and out-phase density oscillations in the two layers. While the high-frequency, in-phase mode, retains the dispersion law of a 2D system $\omega_S \sim \sqrt{q}$, the low-frequency one propagates like an acoustic excitation with $\omega_A \sim q$.[2] When tunneling between the two layers is considered, the out-of-phase mode acquires a gap, leading to a dispersion given by $\omega_- \sim (\Delta^2 + C_1 q + C_2 q^2)^{1/2}$, where Δ defines the plasmon gap, C_1 and C_2 being positive constants related to the material properties.[3] We note here that these results were derived in the cases of 2D or quasi-2D GaAs-type electronic systems, whose energy spectrum in the presence of a magnetic field consists of unpolarized Landau-level minibands, separated by the cyclotron energy $\hbar eB/m^*$.

In the present report, we focus on the study of collective excitations in quasi-2D systems that exhibits a strong spin polarization. This situation is common to dilute magnetic semiconductor heterostructures (quantum wells), whose energy spectrum in a magnetic field is dominated by a large Zeeman splitting, proportional with the gyromagnetic factor γ, sometimes up to a hundred times the band value. Each spin-split subband has a fine structure of Landau levels, but quantum effects associated with the electron orbital motion are weak. A semiclassical description of the electron energy levels in the parabolic approximation is appropriate.

We consider a bilayer structure, formed by two quantum wells of width L, situated in the xy plane separated by a distance d in the z direction. We assume that $L \ll d$, such that inside each layer, the electrons are considered a 2D gas with n electrons per unit area versus a positive background to assure charge neutrality. A constant magnetic field $\vec{B}$ is applied along the $\hat{z}$ axis, such that in equilibrium, the electron population in each of the two wells, assumed identical, is described by an initial spin polarization $0 \leqslant |\xi| \leqslant 1$, where $\xi = (n_\uparrow - n_\downarrow)/(n_\uparrow + n_\downarrow)$, with $n_\uparrow$ and $n_\downarrow$ the spin-up (spin parallel to the applied dc magnetic field) and spin-down (spin antiparallel to the applied dc magnetic field) electron densities, respectively. Any degree of polarization $-1 < \xi < 1$ can be obtained by varying the applied magnetic field.[4,5] The two electron systems are decoupled, no tunneling being considered.

The dielectric and magnetic response functions to an electromagnetic perturbation, whose Fourier transform is described by a momentum $\vec{q}$ and frequency ω dependent electric potential, $\varphi(\vec{q},\omega)$ and magnetic induction $\vec{b}(\vec{q},\omega)$, are obtained by an equation-of-motion method, in a self-consistent approach that include terms beyond the RPA. The exchange and correlation interactions between electrons are considered by generalizing Kukkonen and Overhauser[6] to include spin-dependent local-field corrections $G_\sigma^{\pm}(\vec{q},\omega)$. This idea was first proposed by Yi and Quinn,[7] who analyzed the effect of the polarization factor ξ on the characteristic frequencies of collective modes in a three-dimensional system.

The single-particle Hamiltonian $\mathcal{H}_\sigma$ describes self consistently the effect of the external perturbation and of the induced spin and charge fluctuations. An electron of spin $\vec{\sigma}$ in the ith layer is affected by charge-density fluctuations in both layers $\Delta n^{(i)}$ $(i=1,2)$, and the spin-density fluctuation $\Delta s^{(i)}$ in the same layer. This distinction stresses the long-range action of the Coulomb interaction, whereas the magnetic interaction is assumed to be strong enough only between spins in the same layer. The Fourier component of $\mathcal{H}_\sigma$ can be written as $\mathcal{H}(\vec{q},\omega) = \mathcal{H}_0 + \mathcal{H}_I^\sigma(\vec{q},\omega)$, where $\mathcal{H}_0$ is the equilibrium Hamiltonian and $\mathcal{H}_I^\sigma$ is generated by the spin-dependent self-consistent effective perturbation. This semiclassical description of the electron system is justified by the negligible Landau splitting of the energy levels. The effective perturbation Hamiltonian is

$$\mathcal{H}_I^{\sigma(i)}=\gamma\vec{\sigma}\cdot\vec{b}-e\varphi+v(q)[(1-G_\sigma^+)\Delta n^{(i)}-G_\sigma^-\vec{\sigma}\Delta\vec{s}^{(i)}]+\frac{4\pi\gamma^2}{L}\vec{\sigma}\Delta\vec{s}^{(i)}+v(q)F(q)\times[(1-G_\sigma^+)\Delta n^{(j)}-G_\sigma^-\vec{\sigma}\Delta\vec{s}^{(j)}], \quad (1)$$

where for simplicity the $\vec{\sigma}$ and ω dependence of the local-field parameters, fluctuations, and disturbance is not displayed. The two indices i and j label the two layers of those of the considered system. The others parameters entering Eq. (1) are the Fourier transform of the bare Coulomb interaction, $v(q)=2\pi e^2/\varepsilon q$, the form factor associated with the Coulomb interaction between the two layers, $F(q)=\exp[-qd]$, and the width of the two layers L. (ϵ is the dielectric constant.) These values of the Coulomb interaction are those of a pure 2D system, true in our case only under the assumption that the width of a quantum well is much smaller than the distance between layers. In general, one needs to calculate the expression of the form factors that result from the overlap of the electronic wave functions along the $\hat{z}$ axis. Therefore, the results presented in this paper, in which it is assumed that all electrons have the same subband index, correspond to intrasubband excitations. Intersubband excitations will be analyzed elsewhere.[8]

The local-field corrections $G_\sigma^\pm(\vec{q},\omega)$ can be expressed in terms of correlation and exchange field corrections $G_\sigma^c(\vec{q},\omega)$, as $G_\sigma^\pm(\vec{q},\omega)=G_\sigma^x(\vec{q},\omega)\pm G_\sigma^c(\vec{q},\omega)$. The exact form of the exchange and correlation local-field corrections in a spin-polarized electron system is still an open question. Asymptotic values for large and small wave-numbers q for $G_\sigma^\pm(\vec{q},\omega)$ were obtained in Ref. 9 in the three-dimensional case and in Ref. 10 for the two-dimensional one, by using the equation of motion method.[11–13] In general, the local-field corrections are a function of the polarization factor ξ and the two-particle correlation function at the origin $g(0)$. The asymptotic values for two- and three-dimensional cases are listed in Ref. 10.

The applied electric potential $\varphi(\vec{q},\omega)$ and $b_z(\vec{q},\omega)$ from Eq. (1) generate fluctuations in the number of spins whose orientation remains parallel to the initial polarization direction, along $\hat{z}$. Since the up and down spins oscillate independently, coupled spin- and charge-density excitations, SDE and CDE, respectively, will result.[7] In the first order of perturbation theory, the induced fluctuations are proportional with the effective interaction potential, where the coefficient of proportionality is the appropriate polarization function, $\Pi_{\sigma,\sigma'}$:

$$\Pi_{\sigma\sigma'}(\vec{q},\omega)=\frac{1}{A}\sum_{\vec{k}}\frac{n_{\vec{k}-\vec{q}/2,\sigma'}-n_{\vec{k}+\vec{q}/2,\sigma}}{\hbar\omega-\varepsilon_{\vec{k}+\vec{q}/2,\sigma}+\varepsilon_{\vec{k}-\vec{q}/2,\sigma'}}. \quad (2)$$

(A is the area of the quasi-two-dimensional layer.) Since we assume that all the many-body effects are incorporated in the local-field correction, Eq. (2) is the polarization of the noninteracting electron system. Therefore, $\varepsilon(\vec{k})=\hbar^2\vec{k}^2/2m^*+\gamma B\mathrm{sgn}(\sigma)$ is the equilibrium energy in the initial dc magnetic field, and $n_{\vec{k},\sigma}$ is the usual Fermi distribution function for a quasiparticle with momentum $\vec{k}$ and spin projection σ along the z axis. [The function sgn(σ) is 1 for spin up and -1 for spin down.] Therefore, under the spin-dependent effective potential,

$$V_\sigma^{(i)}=\gamma b_z\mathrm{sgn}\sigma-e\varphi+v(q)[(1-G_\sigma^+)\Delta n^{(i)}-G_{L,\sigma}^-\Delta s^{(i)}\mathrm{sgn}\sigma]+\frac{4\pi\gamma^2}{L}\Delta s^{(i)}\mathrm{sgn}\sigma+v(q)F(q)\times[(1-G_\sigma^+)\Delta n^{(j)}-G_{L,\sigma}^-\Delta s^{(j)}\mathrm{sgn}\sigma], \quad (3)$$

linear fluctuations are $\Delta n_\sigma^{(i)}(\vec{q},\omega)=\Pi_{\sigma\sigma}V_\sigma^{(i)}(\vec{q},\omega)$.

Due to the spin-response anisotropy the local-field correction factors in Eq. (3) are direction dependent: $G_{L,\sigma}^-(\vec{q},\omega)$ being associated to the longitudinal variations, whereas $G_{T,\sigma}^-(\vec{q},\omega)$ to the transverse ones.

In addition to the longitudinal CDE and SDE modes, the interaction of the transverse components of the magnetic field with the electron spin in Eq. (1) generates spin-flip fluctuations or spin waves. Using the usual decomposition of the spin operator, these modes are driven by $(1/2)\sigma_+b_++(1/2)\sigma_-b_-$, where $b_+=b_x+ib_y$ and $b_-=b_x-ib_y$. The effective down-up potential,

$$V_+^{(i)}=\tfrac{1}{2}\gamma b_+-\tfrac{1}{2}v(q)G_{T,\uparrow}^-\Delta s_+^{(i)}-\tfrac{1}{2}v(q)F(q)G_{T,\uparrow}^-\Delta s_+^{(j)}, \quad (4)$$

induces $\Delta n_+^{(i)}(\vec{q},\omega)=\Pi_{\downarrow\uparrow}V_+^{(i)}(\vec{q},\omega)$, while its up-down correspondent

$$V_-^{(i)}=\tfrac{1}{2}\gamma b_--\tfrac{1}{2}v(q)G_{T,\downarrow}^-\Delta s_-^{(i)}-\tfrac{1}{2}v(q)F(q)G_{T,\downarrow}^-\Delta s_-^{(j)}, \quad (5)$$

triggers $\Delta n_-^{(i)}(\vec{q},\omega)=\Pi_{\uparrow\downarrow}V_-^{(i)}(\vec{q},\omega)$.

The collective excitations of the system are obtained for those values of the frequency at which the oscillations are maintained even in the absence of the perturbative field, i.e., the matrix of the susceptibility functions has a null determinant. The transverse modes can be decoupled immediately and their characteristic equation is

$$\left[1+\frac{1}{2}\Pi_{\sigma\sigma'}\left(v(q)f(q)G_{T,\sigma}^--\frac{4\pi\gamma^2}{L}\right)\right]\times\left[1+\frac{1}{2}\Pi_{\sigma\sigma'}\left(v(q)g(q)G_{T,\sigma}^--\frac{4\pi\gamma^2}{L}\right)\right]=0, \quad (6)$$

where $f(q)=1+F(q)$, $g(q)=1-F(q)$, and (σ,σ') can be any of the combinations $(\uparrow,\downarrow)$ or $(\downarrow,\uparrow)$. For a given orientation of the polarization field $\vec{B}$ there are four characteristic frequencies for the possible spin-flip transverse modes. For example, by using the long wavelength limit for the polarization function $\Pi_{\uparrow\downarrow}(\vec{q},\omega)$ (associated with down-up processes), $\Pi_{\uparrow\downarrow}(\vec{q},\omega)\cong n\zeta/\omega-2\gamma B$, we obtain

$$\Omega_{T1}=2\gamma B+\frac{1}{2}n\zeta\left(\frac{4\pi\gamma^2}{L}-v(q)f(q)G^-_{T,\downarrow}\right),$$

$$\Omega_{T2}=2\gamma B+\frac{1}{2}n\zeta\left(\frac{4\pi\gamma^2}{L}-v(q)g(q)G^-_{T,\downarrow}\right). \quad (7)$$

The proportionality of $\Omega_{T2}-\Omega_{T1}$ with $n\zeta e^{-qd}v(q)\times G^-_{T,\downarrow}(\vec{q},\omega)$ allows an independent verification of the transverse local-field correction, $G^-_{T,\downarrow}(\vec{q},\omega)$, by comparison with experimental data.

The characteristic frequencies of the longitudinal modes are the solutions of the following equation:

$$\{1-\beta(\Pi_{\uparrow\uparrow}+\Pi_{\downarrow\downarrow})-g(q)v(q)[\Pi_{\uparrow\uparrow}(1-G^+_{\uparrow}-G^-_{L,\uparrow})+\Pi_{\downarrow\downarrow}(1-G^+_{\downarrow}-G^-_{L,\downarrow})]+2\beta g(q)v(q)\Pi_{\uparrow\uparrow}\Pi_{\downarrow\downarrow}(2-G^+_{\uparrow}-G^+_{\downarrow})$$
$$-2g^2(q)v^2(q)\Pi_{\uparrow\uparrow}\Pi_{\downarrow\downarrow}[G^-_{L,\uparrow}(1-G^+_{\downarrow})+G^-_{L,\downarrow}(1-G^+_{\uparrow})]\}\times\{1-\beta(\Pi_{\uparrow\uparrow}+\Pi_{\downarrow\downarrow})-f(q)v(q)[\Pi_{\uparrow\uparrow}(1-G^+_{\uparrow}-G^-_{L,\uparrow})$$
$$+\Pi_{\downarrow\downarrow}(1-G^+_{\downarrow}-G^-_{L,\downarrow})]+2\beta f(q)v(q)\Pi_{\uparrow\uparrow}\Pi_{\downarrow\downarrow}(2-G^+_{\uparrow}-G^+_{\downarrow})-2f^2(q)v^2(q)\Pi_{\uparrow\uparrow}\Pi_{\downarrow\downarrow}[G^-_{L,\uparrow}(1-G^+_{\downarrow})$$
$$+G^-_{L,\downarrow}(1-G^+_{\uparrow})]\}=0, \quad (8)$$

where $\beta=4\pi\gamma^2/L$. Note that in the RPA, when $G^{\pm}_{\sigma}(\vec{q},\omega)=0$, and in the absence of the self-consistent magnetization, $\beta=0$, Eq. (8) generates the excitation frequencies for the CDE of a bilayer system in the absence of a spin imbalance ($\xi=0$).[2] The solutions of Eq. (8) can be written in terms of the 2D spin-dependent plasma frequencies $\omega^2_{\sigma}=2\pi e^2 n_{\sigma}q/\varepsilon m^*_{\sigma}$, where $n_{\sigma}=n[1+\xi\mathrm{sgn}(\sigma)]/2$ and $m^*_{\sigma}=m^*/\sqrt{1+\xi\mathrm{sgn}(\sigma)}$.[9] With $\alpha=(2\varepsilon\gamma^2/Le^2)q$, the two charge-density modes are obtained to be excited at

$$\Omega_1^2=\alpha(\omega^2_{\uparrow}+\omega^2_{\downarrow})+f(q)[\omega^2_{\uparrow}(1-G^+_{\uparrow}-G^-_{L,\uparrow})+\omega^2_{\downarrow}(1-G^+_{\downarrow}-G^-_{L,\downarrow})]$$
$$-\frac{2\omega^2_{\uparrow}\omega^2_{\downarrow}f(q)\{\alpha(2-G^+_{\uparrow}-G^+_{\downarrow})-f(q)[G^-_{L,\uparrow}(1-G^+_{\downarrow})+G^-_{L,\downarrow}(1-G^+_{\uparrow})]\}}{\alpha(\omega^2_{\uparrow}+\omega^2_{\downarrow})+f(q)[\omega^2_{\uparrow}(1-G^+_{\uparrow}-G^-_{L,\uparrow})+\omega^2_{\downarrow}(1-G^+_{\downarrow}-G^-_{L,\downarrow})]}. \quad (9)$$

$$\Omega_2^2=\alpha(\omega^2_{\uparrow}+\omega^2_{\downarrow})+g(q)[\omega^2_{\uparrow}(1-G^+_{\uparrow}-G^-_{L,\uparrow})+\omega^2_{\downarrow}(1-G^+_{\downarrow}-G^-_{L,\downarrow})]$$
$$-\frac{2\omega^2_{\uparrow}\omega^2_{\downarrow}g(q)\{\alpha(2-G^+_{\uparrow}-G^+_{\downarrow})-g(q)[G^-_{L,\uparrow}(1-G^+_{\downarrow})+G^-_{L,\downarrow}(1-G^+_{\uparrow})]\}}{\alpha(\omega^2_{\uparrow}+\omega^2_{\downarrow})+g(q)[\omega^2_{\uparrow}(1-G^+_{\uparrow}-G^-_{L,\uparrow})+\omega^2_{\downarrow}(1-G^+_{\downarrow}-G^-_{L,\downarrow})]}. \quad (10)$$

By using the asymptotic expression of the polarization function in the long-wavelength limit[2] $\Pi_{\sigma\sigma}=n[1+\xi\mathrm{sgn}(\sigma)]q^2/2m^*\omega^2$, it is then easy to surmise from Eqs. (9) and (10) that in the long-wavelength limit Ω describes the classic propagation of a 2D plasmon $\Omega_1\sim\sqrt{q}$, while Ω_2 corresponds to the acoustic plasmon associated with the out-of-phase charge oscillations described in Ref. 2, with the specification that our calculation includes also the spin corrections.

On account of the spin coupling that generate the self-consistent magnetization, two longitudinal spin-dependent modes arise that describe the in-phase and out-of-phase magnetization fluctuations in the two layers, respectively,

$$\Omega_3^2=\frac{2\omega^2_{\uparrow}\omega^2_{\downarrow}\{\alpha f(q)(2-G^+_{\uparrow}-G^+_{\downarrow})-f^2(q)[G^-_{L,\uparrow}(1-G^+_{\downarrow})+G^-_{L,\downarrow}(1-G^+_{\uparrow})]\}}{\alpha(\omega^2_{\uparrow}+\omega^2_{\downarrow})+f(q)[\omega^2_{\uparrow}(1-G^+_{\uparrow}-G^-_{L,\uparrow})+\omega^2_{\downarrow}(1-G^+_{\downarrow}-G^-_{L,\downarrow})]}, \quad (11)$$

$$\Omega_4^2=\frac{2\omega^2_{\uparrow}\omega^2_{\downarrow}\{\alpha g(q)(2-G^+_{\uparrow}-G^+_{\downarrow})-g^2(q)[G^-_{L,\uparrow}(1-G^+_{\downarrow})+G^-_{L,\downarrow}(1-G^+_{\uparrow})]\}}{\alpha(\omega^2_{\uparrow}+\omega^2_{\downarrow})+g(q)[\omega^2_{\uparrow}(1-G^+_{\uparrow}-G^-_{L,\uparrow})+\omega^2_{\downarrow}(1-G^+_{\downarrow}-G^-_{L,\downarrow})]}. \quad (12)$$

In the absence of the induced magnetization, these modes exist only if $G^-_{L,\sigma}(\vec{q},\omega)<0$, or in terms of exchange and correlation factors, $G^x_{L,\sigma}(\vec{q},\omega)<G^c_{L,\sigma}(\vec{q},\omega)$. Their experimental observation is possible only if they are situated outside the electron-hole continuum. Again, in the long wavelength limit, our analysis involves the asymptotic expression of the polarization function Eq. (2) leading to a regular 2D magnetoplasmon from Eq. (11), and to an acoustic magnetic excitation from Eq. (12).

This latter mode has the lowest frequency of the four modes and will have the strongest contribution to possible decay processes. It is then important to analyze under what circumstances it will exist. If all the local-field correction factors are set to zero (RPA), from Eq. (12) we obtain in the long-wavelength limit:

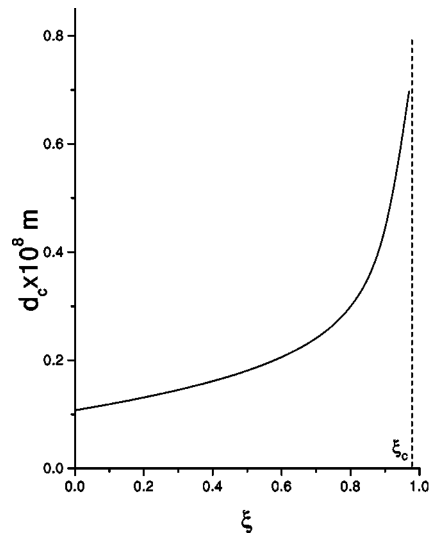

FIG. 1. The dependence of the critical interlayer distance d_c on the polarization factor, ξ, is presented for $L=2\times10^{-8}$ m, $\varepsilon\simeq12$, $\gamma\simeq100\mu_B$, and $m^*\simeq0.1m_e$. For these values, the critical polarization factor $\xi_c\simeq0.99$.

$$\Omega_4^2=v_F^2q^2\left(\frac{\sqrt{1-\xi^2}}{\sqrt{1+\xi}+\sqrt{1-\xi}}\right)\frac{\frac{q_{TF}\varepsilon\gamma^2}{Le^2}}{1+\frac{2\varepsilon\gamma^2}{e^2Ld}}, \tag{13}$$

with $q_{TF}=2m^*e^2/\varepsilon\hbar^2$ the usual Thomas-Fermi wave number. As in the case of charge excitations, this mode exists only for a finite distance between the layers. Moreover, this distance has to be greater than a critical value for the collective excitation to fall outside the electron-hole continuum. This results from the condition that the the group velocity of the collective mode is greater than the higher of the two Fermi velocities, $v_{F\sigma}=v_F[1+\xi\mathrm{sgn}(\sigma)]$. (Note that in order to include such effects, one should consider the spin dependence of both the Fermi momentum and the effective electron mass.) Therefore, for $\xi>0$, when $v_{F\uparrow}>v_{F\downarrow}$ the interlayer distance should satisfy:

$$d>d_c=\frac{\frac{2\varepsilon\gamma^2}{Le^2}}{\frac{2\varepsilon\gamma^2}{Le^2}q_{TF}\frac{\sqrt{1-\xi}}{1+\xi+\sqrt{1-\xi^2}}-1}. \tag{14}$$

The critical distance between the two layers thus determined is spin dependent, via the polarization factor ξ. The denominator of Eq. (14) is required to be positive, which imposes that

$$\frac{2\varepsilon\gamma^2}{Le^2}q_{TF}\frac{\sqrt{1-\xi}}{1+\xi+\sqrt{1-\xi^2}}>1. \tag{15}$$

This condition limits dramatically the number of instances in which the effect discussed here can be observed. For a quantum well of a typical width of 100 Å, Eq. (15) is satisfied only for large values of the gyromagnetic factor γ. This is why we suggest that the existence of an acoustic magnetoplasmon can be observed only in dilute magnetic semiconductor structures whose γ can be as high as 100 μ_B. For a given layer thickness, L, Eq. (15) is satisfied only up to a critical polarization value, ξ_c. For well widths of the order of hundreds of angstroms, the average size of a quantum well, ξ_c is close to one, so the magnetoacustic plasmon we predict should be observed for different values of the polarization. For larger widths, the polarization parameter is limited in agreement with Eq. (15). In Fig. 1, we analyze the dependence of the critical interlayer distance, d_c as function of the polarization factor ξ for a layer thickness $L=2\times10^{-8}$ m. Close to ξ_c, the critical distance increases abruptly signaling the softening of the magneto-acoustic collective mode, which can not be observed any more.

In conclusion, we present an analysis of the possible transverse and longitudinal collective modes in a dilute semiconductor bilayer structure. Transverse spin fluctuations are associated with spin waves as expected and the difference of their excitation frequencies can be directly associated with the transverse local-field correction. We show that four collective excitations can result from the coupling of charge and spin fluctuations: in-phase and out-of-phase charge density excitations, and in-phase and out-of phase spin-density excitations. The SDE and CDE coupling is mediated by the self-induced magnetization. The out of phase modes are acoustic, both for charge and spin. The out-of-phase SDE is described by an acoustic dispersion, which if exists outside the electron-hole continuum represents the lowest frequency possible for a collective excitation in this system. Based on our results obtained in the RPA, we argue that it will be possible to observe the acoustic magnetoplasmon in a material with high γ, such as a dilute magnetic semiconductor.

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