Effect of Pauli blockade on spin-dependent diffusion in degenerate semiconductors

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In p+ GaAs thin films, diffusion of degenerate photoelectrons is shown to strongly depend on spin. This is a direct consequence of the Pauli principle which forbids elementary scattering processes if the scattered electronic state is already occupied by an electron of the same spin. Spin-polarized (40%) photoelectrons are created at 15 K by a tightly-focused circularly-polarized light excitation and the spin-dependent diffusion profile is monitored by imaging the luminescence intensity and polarization as a function of electronic temperature and excitation power. For nondegenerate concentrations the spin polarization simply decreases with distance from the excitation spot as determined by the spin lifetime and the essentially spin-independent diffusion constant. In sharp contrast, in degenerate conditions, Pauli blockade produces a spin filter effect which results in a depletion of majority spin electrons at the center and the appearance of a polarization maximum at a distance of about 2 \( \mu m \) from the center. Spin-dependent diffusion is directly evidenced from the observed dependence of the luminescence intensity profile on the electronic spin polarization. The results are in quantitative agreement with a resolution of the coupled charge and spin diffusion equations, using no adjustable parameter, and imply a relative difference of about 55% between the diffusion constants of majority and minority spin electrons and of 20% between their mobilities.

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I. INTRODUCTION

Investigation of spin-dependent transport properties in semiconductors has been a very active field of research for more than 40 years. A number of investigations considered spin-dependent photoconductivity [1, 2], induced by spin-dependent recombination [3], or spin-dependent Hall effect [4-6]. Further, the possibility to use spin degrees of freedom to transport information[7] has triggered very active recent research on the various possible mechanisms of coupling between spin and charge transport. In this context, different coupling mechanisms have been considered, such as the spin Coulomb drag explored by Flensberg et al. [8] and by D’Amico and Vignale [9], by which electron-electron interactions introduce an intrinsic source of friction for spin currents, reducing the spin diffusion constant relative to the charge diffusion constant and resulting in nonzero off-diagonal elements of the diffusion matrix. This spin Coulomb drag was recently observed by Weber et al. [10]. Another example is the persistent spin helix observed in quantum wells [11].

Spin-dependent transport should be strongly amplified in conditions where the electron gas is degenerate. As a consequence of the Pauli exclusion principle that dictates that, elementary scattering processes are in such a case less likely. The diffusion constant is then larger than that predicted by the Einstein relation [12]. In the case of a spin-polarized electron gas, this implies that the diffusion constant of majority spin electrons is larger than that of minority spin electrons. This Pauli blockade was observed in a localized electronic system consisting of a single quantum dot [13], or in Si/SiGe double quantum dots [14], but not yet in transport in a delocalized electron gas. While degeneracy has been included in recent theoretical calculations of spin-polarized transport [15-16], and in experimental studies such as the spin-flip blocking in bulk GaAs [17], its explicit implications for the diffusion of spin polarization and spin-charge coupling have not been explored.

Here we demonstrate explicitly, theoretically and experimentally, the existence of a new spin-charge coupling mechanism in a degenerate photoelectrons gas. This mechanism results in strongly spin-dependent diffusive and drift transport and yields a spatial spin-filter effect. We use a spatially-resolved polarized microluminescence technique. Spin-polarized carriers are created at 15K by a tightly-focussed circularly-polarized laser and the spatial distribution of charge and spin are investigated using polarized imaging. We find an increase of spin polarization after creation during diffusive transport, in strong contrast with the usual decrease due to spin-lattice re-
laxation. A second effect, of second order in the polarization, is that the charge diffusion constant of the photoelectrons depends on their spin polarization, as seen from the dependence of the intensity profile on polarization of the exciting light. The results are quantitatively interpreted using parameter values determined independently. As a result, the diffusion constant of majority spin carriers is found larger than that of minority spin carriers. In the same way and for the same physical reasons, the mobility of majority spin carriers is found larger than that of minority spin carriers by about 20%.

II. THEORY

A. Spin dependent diffusion induced by Pauli blockade

In a non-degenerate semiconductor, the diffusion constant is given by the Einstein relation

\[ D_0 = \mu_0 k_B T_e / q \]  

where \( k_B \) is the Boltzmann’s constant, \( q \) the absolute value of the electron charge, \( T_e \) the temperature of the electron gas and \( \mu_0 \) is the mobility, given by

\[ \mu_0 = q \tau_m / m^* \]  

where \( m^* \) is the effective mass in the conduction band and \( \tau_m \) is the characteristic time of the scattering processes, such as electron-phonon scattering, minority-majority carrier scattering, or electron-ionized impurity scattering. As illustrated in the inset of the bottom panel of Fig. 1 in a non-degenerate regime, the crystal momentum \( \vec{k}_i \) of a Bloch electron in the conduction band can change due to the presence (for example) of an ionized impurity, finishing with a wavevector \( \vec{k}_f \) at the end of the scattering process. In contrast, in the degenerate case, illustrated in the inset of the top panel of Fig. 1, the Pauli exclusion principle tells us that the corresponding elementary scattering process can be forbidden if the final state after scattering is already occupied by an electron of the same spin, that results in an increase of the diffusion constant.

Here, we consider that circularly-polarized light excitation of a p-type semiconductor creates a degenerate spin-polarized density of photoelectrons in the conduction band. The thermal distributions of the two species of photoelectrons, with spin aligned parallel or anti-parallel to the direction of light excitation, are described by two distinct quasi-Fermi energies \( E_{F\pm} \), related to the concentrations \( n_\pm \) of the corresponding electrons by

\[ n_\pm = \frac{2}{\sqrt{\pi}} N_e^s \mathcal{F}_{1/2}(E_{F\pm}) \]  

where \( E_{F\pm} = 0 \) at the conduction band edge, and the Fermi integral is defined by

\[ \mathcal{F}_n(E_{F\pm}) = \int_0^\infty x^n f_\pm(x) dx \]  

with \( f_\pm(x) \) the Fermi-Dirac distribution at an energy \( k_B T_e x \) and Fermi level \( E_{F\pm} \). Here the weak dependence of the electron effective mass on temperature is neglected and the effective spin density of states in the conduction band is

\[ N_e^s = N_e / 2 = 2.35 \times 10^{17} (T_e/300)^{3/2} \text{ cm}^{-3} \]  

FIG. 1: Spin-resolved photoelectron concentration dependence of the reduced diffusion constant \( \nu [\text{Eq. (9)}] \), of the reduced mobility \( \zeta [\text{Eq. (5)}] \), of the proportionality coefficient \( \xi \) between \( D \) and \( \mu \) [Eq. (6), squares] and of the Fermi energy for \( T_e = 50 \text{ K} \) (top panel), and for \( T_e = 150 \text{ K} \) (bottom panel). Shown by dots are the values of \( n_\pm \) and \( n_\pm \) calculated with \( n = 10^{17} \text{ cm}^{-3} \) and \( \mathcal{F} = 30 \% \). Comparison with the spin-resolved density of states \( N_e / 2 \), given by Eq. (9), shows that the top (bottom) panel corresponds to degenerate (non degenerate) conditions. As illustrated in the insets, elementary scattering processes are likely to be forbidden in degenerate conditions, because of the Pauli principle. As a result, the reduced mobility \( \zeta \) and the proportionality coefficient \( \xi \) between \( D \) and \( \mu \) depend more on concentration in degenerate than in non-degenerate conditions.
In a degenerate gas, Eq. \[1\] should be replaced by \[12\]:

\[D_\pm = \left(\frac{\mu_\pm k_B T_e}{q}\right) \frac{2}{\mathcal{F}_{\pm 1/2}(E_{F_\pm})} = \xi(n_\pm)\frac{\mu_\pm k_B T_e}{q} \] \[6\]

where the mobility \(\mu_0\) appearing in Eq. \[2\] is replaced by \(\mu_\pm\). The latter quantity is now spin-dependent via the mean relaxation time \[12\]

\[\tau_m = \int_0^{\infty} \frac{3^2}{\mathcal{F}_2(E_{F_\pm})} \frac{\partial f_\pm(\varepsilon)}{\partial \varepsilon} \tau(\varepsilon) d\varepsilon / \int_0^{\infty} \frac{3^2}{\mathcal{F}_2(E_{F_\pm})} \frac{\partial f_\pm(\varepsilon)}{\partial \varepsilon} d\varepsilon \] \[7\]

While the dominant photoelectron scattering mechanism at low temperature is not completely known \[20\], it is reasonable to assume that ionized impurity scattering plays a key role, in which case \(\tau_m \propto 3^2/2\) \[12\]. One then finds

\[\mu_\pm = \mu_0 \sqrt{\pi} \mathcal{F}_2(E_{F_\pm}) 4 \mathcal{F}_{-1/2}(E_{F_\pm}) = \mu_0 \zeta(n_\pm) \] \[8\]

Using Eqs. \[6\] and \[8\], one obtains

\[D_\pm = \mu_0 \left(\frac{k_B T_e}{q}\right) \frac{\sqrt{\pi}}{2} \mathcal{F}_2(E_{F_\pm}) = D_0 \nu(n_\pm) \] \[9\]

where \(D_0\) is given by Eq. \[11\], and \(\nu(n_\pm) = \xi(n_\pm) \zeta(n_\pm)\).

The latter function is larger than unity as soon as the spin electron concentration \(n_\pm\) is comparable with \(N_e^s\), given by Eq. \[4\]. In the strongly degenerate regime \((n_\pm \gg N_e^s)\), one finds \(D_\pm = (\sqrt{\pi}/12)D_0 \left((E_{F_\pm} - E_e)/k_B T_e\right)^{5/2} \propto n_e^{5/3}\). For highly p-doped GaAs, Fig. \[4\] shows the concentration dependences of \(D_\pm/D_0\) and of the two functions \(\xi\) and \(\zeta\) for two different temperatures. Note that the concentration dependences of \(\xi\) and \(\zeta\) are found similar. Also shown are the concentration dependences of the Fermi energies.

A typical electron gas of density \(n_+ + n_- = 10^{17}\) cm\(^{-3}\) with a spin polarization of \(\mathcal{P} = (n_+ - n_-)/n = 30\%\) is represented by the two photoelectron densities shown in Fig. \[4\] by vertical dashed lines. At a temperature of \(T_e = 150\) K both concentrations are smaller than the effective spin-resolved density of states, and \(D_+/D_-\) is equal to 1.15. In contrast, at a lower temperature of \(T_e = 50\) K, both densities are larger than \(N_e/2\) and \(D_+/D_- \approx 1.78\). This larger value illustrates the fact that, at a fixed value of \(n_\pm\), the effect of degeneracy on diffusion increases as the electronic temperature is lowered due to the variation of \(N_e\) with temperature. It is concluded that Pauli blocking can lead to a significant spin dependence of the diffusion constant. The purpose of this work is to demonstrate this effect and to explore its implications for spin transport.

### B. Diffusion matrix

In the present section, we give the expression of the diffusion matrix taking account of the Pauli blockade and of the possible couplings between up and down spins, due to which a spatial gradient \(\nabla n_\alpha\) of the concentration of a given spin species induces a current of the other species \(n_\beta\). Since this spatial gradient induces an electric field \(-q(\partial E_{F_\alpha}/\partial n_\alpha)\nabla n_\alpha\) acting on the corresponding electrons, the diffusion of the photoelectrons is described by a matrix \(D_{\alpha\beta}\), that can be written \[13\] as \(q^2 D_{\alpha\beta} = \sum_\gamma \sigma_{\alpha\gamma} S_{\gamma\beta}\), where \(\sigma\) is the conductivity matrix and the spin stiffness matrix is defined by \(S_{\alpha\beta} = \partial E_{F_\alpha}/\partial n_\beta\).

As shown elsewhere, we consider here two distinct coupling mechanisms \[13\][10]. The first one is due to the dependence of the bandgap on the total electron concentration. This bandgap renormalization causes electrons of a given spin orientation to feel the electronic concentration of the opposite orientation. The other mechanism which leads to charge-spin coupling is the so-called spin-drag, that depends on the relative efficiency \(1/\tau_{ee}\) of electron-electron collisions with opposite spin, with respect to the efficiency of all other scattering mechanisms \[21\]. Since as shown in Eq. \[7\] the latter efficiency is also spin dependent, we shall use an effective spin-averaged time given by

\[\tau_m = \frac{n_+}{n} \tau_m - + \frac{n_-}{n} \tau_m + \] \[10\]

As a result, the effective spin diffusion constant, defined by

\[D_{\alpha\beta} = D_0 / (1 + \tau_m/\tau_{ee}) \] \[11\]

can be smaller than \(D_0\) due to spin-drag. The general result is then written

\[D_{\alpha\beta} = \frac{n_\alpha \mu_\alpha}{q(1 + \tau_m/\tau_{ee})} \left[1 + \frac{\tau_{ee}}{\tau_{ee}} n \frac{\partial E_{F_\alpha}}{\partial n_\beta} + \frac{\tau_{ee}}{\tau_{ee}} n \frac{\partial E_{F_\alpha}}{\partial n_\beta}\right] \] \[12\]

with \(n_{\alpha\beta}\) the electron density with spin opposite to \(\alpha\). Since \(n_\alpha \partial E_{F_\alpha}/\partial n_\alpha = k_B T_e \xi(n_\alpha)\), Pauli blockade contributes to the terms of the form \(\partial E_{F_\alpha}/\partial n_\alpha\) in the diffusion matrix, while renormalization effects give rise to non-zero values of \(\partial E_{F_\alpha}/\partial n_\alpha\). Assuming for simplicity that bandgap renormalization effects are negligible, the expression for the diffusion matrix given by Eq. \[12\] becomes

\[D = \frac{D_0}{1 + \tau_{ee}} \left(\frac{\nu(n_+)}{\tau_{ee}} n \frac{\nu(n_-)}{\nu(n)} \right) \] \[13\]

where, to first order in \(\mathcal{P}\), one has

\[\nu(n_\pm) = \nu(n/2)[1 \pm \eta \mathcal{P}] \] \[14\]

where

\[\eta = \frac{n/2}{\nu(n/2)} \frac{d \nu(n/2)}{dn} \] \[15\]
In the strongly degenerate regime, $\eta = 5/3$.

C. Diffusion equations

Because of electrostatic effects which induce ambipolar diffusion [22], diffusion of charge and spin is coupled to the diffusion of photocreated holes. The corresponding diffusion equations are

$$g_+ (\vec{r}) - K(N_A + \delta p)n_+ - (n_+ - n_-)/2T_1 + \vec{\nabla} \cdot (\mu_+ n_+ \vec{E} + D_{++} \vec{\nabla} n_+ + D_{+-} \vec{\nabla} n_-) = 0$$

$$g_-(\vec{r}) = K(N_A + \delta p)n_- + (n_+ - n_-)/2T_1 + \vec{\nabla} \cdot (\mu_- n_- \vec{E} + D_{--} \vec{\nabla} n_- + D_{-+} \vec{\nabla} n_+) = 0$$

$$g(\vec{r}) = K(N_A + \delta p)(n_+ + n_-) + \vec{\nabla} \cdot (\mu_h (\delta p + N_A) \vec{E} + D_h \vec{\nabla} \delta p) = 0$$

Here, $\delta p$ is the excess of photoholes concentration created by light. $g_\pm (\vec{r})$ is the rate of spin $\pm$ electron-hole pair creation, which may depend on position $\vec{r}$ for a strongly focused excitation, $g = g_+ - g_-$. $T_1$ is the spin-lattice relaxation time, $N_A$ is the acceptor density, $K$ is the bimolecular recombination coefficient so that at low excitation power the bulk minority carrier lifetime is given by $1/\tau = KN_A$. The quantity $\mu_h$ is the hole mobility and $N_A$ is the density of ionized acceptors. Finally, $\vec{E}$ is the internal electric field that satisfies Poisson’s equation: $\vec{\nabla} \cdot \vec{E} = \varepsilon (\delta p - n)/\varepsilon$ where $\varepsilon$ is the permittivity.

Using Eqs. [13] and [14], Eqs. [13] and [17] are then combined in order to eliminate the terms dependent on electric field. This gives coupled diffusion equations implying the electronic charge $\langle n = n_+ + n_- \rangle$ and spin $\langle s = n_+ - n_- \rangle$ densities

$$g - \frac{n}{\tau} + \vec{\nabla} \cdot [D_a \vec{\nabla} n + D_a^s \eta \mathcal{P} \vec{\nabla} s] = 0$$

$$g\mathcal{P}_h - \frac{s}{\tau_s} + \vec{\nabla} \cdot \left[ D_a^s \vec{\nabla} s + D_s \left[ \eta \mathcal{P} + \frac{D_a^s \tau_m}{D_0 \tau_{ee}} \right] \vec{\nabla} n \right] = 0$$

where $\tau_s = (1/T_1 + 1/\tau)^{-1}$ is the spin lifetime and

$$D_a^s = D_a \nu (n/2)$$

To first order in $\mathcal{P}$, and assuming $n \approx \delta p$ [22], the Coulomb interaction between photoelectrons and photoholes is taken into account via the following ambipolar diffusion constants [22]:

$$D_a = D_0 \nu (n/2) \frac{n[1 + \xi^{-1}(n/2)] + N_A^{-}}{n[1 + \zeta(n/2)\mu_0/\mu_h] + N_A}$$

and

$$D_a^s = D_0 \nu (n/2) \frac{n + N_A^{-}}{n[1 + \zeta(n/2)\mu_0/\mu_h] + N_A}$$

The ratio $\mu_0/\mu_h = m^*/m_e$ is equal to 7.9 [22], where $m_e$ is the effective mass of holes. Note that Eq. (19) now contains a coupling term to the spin-diffusion equation. This term is proportional to $\eta \mathcal{P} = \xi^2 \mathcal{P}$ for $n \gg N_c$, and therefore increases with electron polarization. The latter equation is not modified by spin-drag, since the total charge current is not altered by electron-electron collisions. In contrast, the coupling term in the spin diffusion equation also includes a similar term, proportional to $\eta \mathcal{P}$, with an additional contribution due to spin drag. Considering a sample of thickness $d$ with very large lateral dimensions, the boundary conditions are given at the front surface ($z = 0$) by

$$D_a \partial n/\partial z + \frac{D_a^s \eta \mathcal{P} \partial s/\partial z = S_n}{\partial s/\partial z = S_s}$$

and

$$D_s \partial s/\partial z + D_s [\eta \mathcal{P} + \frac{\tau_m D_a^s}{D_0 \tau_{ee}}] \partial n/\partial z = S_s$$

The corresponding boundary conditions at the back surface of the sample ($z = d$) are simply obtained by replacing in the latter equations $S$ by $-S$.

Note that spin transport can in principle be affected by spin drag, since the corresponding effects appear in the spin diffusion equation [Eq. (20)] as a contribution to the coupling term which is added to Pauli blockade effects. In contrast, the charge diffusion equation only depends on Pauli blockade, so that the most direct verification of the latter effect should be to experimentally investigate the dependence of the photoelectron concentration as a function of their average polarization $\mathcal{P}$. This will reveal a polarization dependence of the spin-averaged diffusion constant, defined by $\langle D \rangle = (n_+ D_{++} + n_- D_{--})/n$, which is found equal to

$$\langle D \rangle = D_a \left( 1 + \eta \mathcal{P}^2 \right)$$

and depends on spin via a second order effect. Qualitatively speaking, when the electrons are polarized, majority electrons are forced to occupy high energy levels due to the Pauli principle, creating thus an internal pressure of the gas that is higher than in the case of an unpolarized gas.
III. EXPERIMENT

A. Principle of the experiment

P-type \( (N_A \approx 10^{18} \text{ cm}^{-3}) \) GaAs films of thickness \( d = 3 \mu \text{m} \) have been grown on a GaAs semi-insulating substrate, with a thin GaInP back layer serving as a confinement layer for the photoelectrons and ensuring a negligible recombination velocity \( S' = 0 \) at the GaAs back surface, as shown in the bottom panel of Fig. 2. The front surface is naturally oxidized.

Charge and spin diffusion have been investigated between 6 \( K \) and 300 \( K \), using a polarized microluminescence technique described elsewhere [24]. A circularly-polarized light excitation at 1.59 eV is focused to a Gaussian spot of Gaussian half width \( \omega = 0.6 \mu \text{m} \) by using a microscope adapted to analyze the polarization of the luminescence. A schematic of this microscope is shown in the top panel of Fig. 2. Since the laser light is linearly-polarized when it enters the microscope, its polarization at the sample can be switched between \( \sigma^+ \), \( \sigma^- \) and \( \pi \) by rotating a half-wave plate with respect to a quarter-wave plate. The photoluminescence (PL) of the sample is analyzed by a second quarter-wave plate followed by a linear polarizer, enabling to collect only the \( \sigma^+ \) or \( \sigma^- \) polarized components of the luminescence, depending on the orientation of the quarter-wave plate. In a first stage, both \( \sigma^+ \) and \( \sigma^- \) polarized light excitations are in turn used to excite the sample, and an image (or spectrum) is taken of the \( \sigma^\parallel \) polarized components of the resulting luminescence emission with the laser being removed by an appropriate filter. The resulting four images, denoted \( \sigma^+ \), \( \sigma^- \), \( \pi^+ \) and \( \pi^- \), are combined to form a sum image \( I_s = (\sigma^+ + \sigma^- + \pi^+ + \pi^-)/2 \) proportional to \( n \) and a difference image \( I_d = (\sigma^+ - \sigma^- + \pi^- - \pi^+)/2 \) proportional to \( s = n_+ - n_- \). By dividing these two images, an image proportional to the electronic polarization is obtained. Their respective sections as a function of radial distance \( r \) to the excitation spot are given by

\[
I_s(r) = A \int_0^d n(r,z) \exp[-\alpha_l z] \, dz
\]

\[
I_d(r) = A |\mathcal{P}_t| \int_0^d s(r,z) \exp[-\alpha_l z] \, dz
\]

and

\[
I_d(r)/I_s(r) = \mathcal{P}(r) \mathcal{P}_t.
\]

Here \( A \) is a constant and \( \alpha_l \approx (3 \mu \text{m})^{-1} \) is the absorption coefficient at the luminescence energy [23]. The quantity \( \mathcal{P}_t \) depends on the matrix elements for recombination and is equal to \( \pm 0.5 \) for \( \sigma^\parallel \) light excitation [25].

A spatially-resolved spectrum of the luminescence can be obtained by using a beam splitter to send some part of the light to the entrance of a multi-mode optical fiber, at the output of which it is analyzed by a spectrometer and
a CCD camera which gives the emission spectrum (see Fig. 2). Depending on the radius of the optical fiber, it is possible to record a spatially averaged spectrum, or a spatially-resolved spectrum in an adjustable point of radius 0.9 μm in the image.

**B. Sample characterization**

In the present section, we estimate the key parameters for the investigation of Pauli blocking. These parameters are charge and spin diffusion constants, conduction electron lifetime and spin-lattice relaxation times. The measured values will be used throughout Sec. (V), so that quantitative interpretation of the experimental results will be performed without any adjustable parameter. Sample characterization was first performed using the photoluminescence imaging technique described in Ref. [26], and peaks at 1.50 eV. From the measured shape of the high energy tail, the value of the photoelectron temperature \( T_e = 50 \) K is obtained. As found in the same figure, the luminescence circular polarization is of the order of 19 % and almost independent of photon energy, so that the mean photoelectron polarization, of 38 %, is close to its maximum value of \( P_i = 50 \) %. The difference can be due to losses during thermalization or due to spin relaxation during the lifetime of the thermalized electrons in the conduction band. In the latter case, one has according to \( I = P_i \times \tau_{s,eff}/\tau_{eff} \), where the times \( \tau_{eff} \) and \( \tau_{s,eff} \) are the effective minority carrier and spin lifetimes, including surface recombination [22]. The time \( \tau_{eff} \) is defined by

\[
\frac{1}{\tau_{eff}} = \frac{1}{\tau} + D_0 \alpha_1^2 \tag{30}
\]

where \( \alpha_1 d \) is the positive solution smaller than \( \pi/2 \) of

\[
\alpha_1 d \tan(\alpha_1 d) = Sd/D_0 \tag{31}
\]

and the time \( \tau_{s,eff} \) is given by a similar equation where \( \tau \) is replaced by \( \tau_s \).

Also shown in the top left panel of Fig. 3 are the spatial cross sections of the sum and difference micro-luminescence images, as defined by Eq. (27) and Eq. (28) respectively, and the excitation laser profile. These curves are in very good agreement with a numerical solution of Eq. (19) and Eq. (20) taking \( \eta = 0 \), \( \nu(\pi/2) = 1 \), \( \tau_m/\tau_{ee} = 0 \) and \( D_a = D^*_a = D_0 \). Using Ref. [22], we find

\[
L_e = \sqrt{D_0 \tau_{s,eff}} \approx 1.45 \pm 0.05 \mu m \tag{32}
\]

Since the above measurements do not enable separate determination of the diffusion constants and of the lifetimes, a time-resolved photoluminescence using a setup described in Ref. [28] was performed. The excitation source was a circularly-polarized mode-locked frequency doubled Ti:Sa laser with a 1.5 ps pulse width and a repetition frequency of 80 MHz. The \( \sigma^+ \) - and the \( \sigma^- \) -polarized components of the luminescence were monitored separately, analyzed by a spectrometer and detected by a streak camera. The time-resolved PL spectrum is close to the one shown in Fig. 3. Shown in curves a and b of the bottom panel of Fig. 4 is the time evolution of the sum and difference quantities, taken at a temperature of 20 K. These two quantities are related to the charge and spin orientation by equations identical to Eq. (24) and (23). The only difference is that, since the diameter of the laser beam (of about 50 μm), is much larger than \( L_e \) and \( L_s \), the charge and spin are homogeneous as a function of distance \( r \) so that diffusion is one dimensional along the \( z \) direction defined in Fig. 2.

Both curves a and b exhibit an exponential decrease between about 300 psec and 1000 psec after the laser pulse with a faster decrease at the beginning. This decrease is independent of power over the whole experimental range, i.e. between 5 µW and 500 µW average power. In order to perform a more detailed analysis, an efficient algorithm was used to solve the following 1D inverse Fredholm integral of the first kind

\[
I(t) = \int_{\tau_{min}}^{\tau_{max}} J(\tau)e^{-t/\tau}d\tau + \varepsilon(t) \tag{34}
\]

where \( I(t) \) represents the sampled measurements, (Here \( I_s \) or \( I_d \) \( J(\tau) \) is the searched distribution of time constants and \( \varepsilon(t) \) represents the random noise. This algorithm is based on an inverse Laplace transform containing a compression of the data obtained with the singular value decomposition of the kernel and a robust optimization method using a Tikhonov regularization. Though the inverse Laplace method is known as an ill-posed problem, the very good signal to noise of the experimental data gives two well-defined peaks. These peaks are at \( \tau_1 = 257 \pm 15 \) psec and \( \tau_2 = 66 \pm 20 \) psec for the sum transient and at \( \tau_{1s} = 213 \pm 15 \) and \( \tau_{2s} = 62 \pm 20 \) psec for the difference one, as shown in the top panel of Fig. 4. No other contribution is present in the whole time range between \( \tau_{min} = 10 \) psec and \( \tau_{max} = 1000 \) psec.

This result is in agreement with the expected presence of modes for a system of photoelectrons undergoing one dimensional diffusion and surface recombination [30]. While complete analysis of the times and amplitudes of these modes will be performed elsewhere, we limit ourselves here to an estimate of the relevant parameters. The longest times \( \tau_1 \) and \( \tau_{1s} \) are equal to \( \tau_{eff} \) given by Eq. (30) and \( \tau_{s,eff} \), respectively. For charge diffusion, this
is the characteristic decay time of the circular polarization luminescence experiment. Since weak, at least up to the maximum power of the transient \[15\] within experimental uncertainty, so that spin-drag is factor of 4 for material of smaller quality (bulk) \[33\]. The For similar doping and temperature, a value larger by a measure of the photoelectron intensity Curve \(a\) of the laser spot larger than charge and spin diffusion lengths. FIG. 4: The bottom panel shows the time-resolved characterization of the sample dynamic properties at 20 K, with a size of the laser spot larger than charge and spin diffusion lengths. Curve \(a\) shows the time evolution of the sum signal, which is a measure of the photoelectron intensity \(n_+ + n_-\), as a function of time after the excitation pulse. Curve \(b\) shows the time evolution of the difference signal, proportional to \(n_+ - n_-\). Curve \(c\) is the ratio of the two signals equal to the degree of circular polarization of the emitted light. Shown in the top panel are inverse Laplace transforms of curves \(a\) and \(b\), as defined by Eq. \[34\]. These spectra exhibit narrow peaks which give the characteristic times of the parts of the transients at long and short delay after the pulse.

gives \(D_0 = L_0^2/\tau_{eff} \approx 82 \pm 10 \text{ cm}^2/\text{s}\). This value can be compared with the room temperature diffusion constant, since the mobility of minority carrier is approximately proportional to \(1/\text{T}\) so that, as seen from Eq. \[11\], the diffusion constant should weakly dependent on temperature \[20\]. Using Ref. \[31\], we estimate that at 300 K, \(D_0 \approx 40 \pm 10 \text{ cm}^2/\text{s}\) \[32\]. The spin diffusion constant is found equal to \(D_s = L_s^2/\tau_{s,eff} \approx 71 \pm 15 \text{ cm}^2/\text{s}\). Examination of Eq. \[11\] allows us to conclude that \(\tau_m \ll \tau_{ee}\) \[15\] within experimental uncertainty; so that spin-drag is weak, at least up to the maximum power of the transient luminescence experiment. Since \(D_0 \approx D_s\), it is seen that the characteristic decay time of the circular polarization is \(T_1\). As shown in curve \(c\) of Fig. 4 this decay is exponential, from which \(T_1 = 1250 \pm 100 \text{ ps}\) is obtained. For similar doping and temperature, a value larger by a factor of 4 for material of smaller quality (bulk) \[33\]. The time of the faster mode is given by

\[
\frac{1}{\tau_2} = \frac{1}{\tau} + D(\alpha_2 + \pi/d)^2 \tag{35}
\]

where \(\alpha_2\) satisfies the equation

\[
(\alpha_2d + \pi)\tan(\alpha_2d) = \alpha_1d\tan(\alpha_1d) = Sd/D \tag{36}
\]

For the difference transient the same equations are also valid, provided \(\tau\) is replaced by the spin lifetime \(\tau_s\). Eqs. \[15\] and \[36\] yield a numerical equation which depends only on \(\tau\), from which we obtain \(\tau \approx 270 \pm 50\) psec. This value is to be compared with the estimated value of the radiative lifetime, of 340 ps, obtained from the value at 300 K of about 5 ns for the same doping \[34\] and applying a power-law dependence of exponent 1.5 as a function of temperature \[30\] between 300K and the actual value of 50 K of the temperature of the electron gas. Finally, Eq. \[31\] gives \(S = 4.6 \times 10^4 \text{ cm/sec}\) although the uncertainty of this determination is very large because of the close values of \(\tau_1\) and \(\tau\). Due to the small value of \(S\) at low temperature, the boundary equations \[21\], \[25\], reduce to \(\partial n/\partial z = 0\) and \(\partial s/\partial z = 0\) at the two interfaces.

IV. RESULTS

A. Spin filter effect

The most spectacular manifestation of the Pauli blockade is on the spatial profile of the luminescence polarization \(P\), as defined by Eq. \[20\]. Shown in Fig. 5 are images of \(P\) taken at 15 K for four different excitation powers, a) 72 nW, b) 0.45 mW, c) 1.89 mW and d) 2.55 mW. As the power is increased, a polarization dip at the center progressively appears. Remarkably, over the first 2 \(\mu\)m, the polarization increases as a function of distance to the excitation spot. Such finding is highly counterintuitive since, under usual diffusion conditions, one expects as shown in Panel a, the polarization to be a monotonic decreasing function of distance \[24\]. While it will be shown below that these results can be explained quantitatively using Eqs. \[19\] and \[20\], qualitative explanation, considering only Pauli blockade, is given in the inset of Fig. 5. This inset shows separate diffusion profiles of majority \(n_+\) and minority \(n_- \) electrons, assuming a strongly-localized light excitation at \(r = 0\) and \(D_{++} > D_{--}\), where \(D_{\pm\pm} = D_0\nu(n_{\pm})\). It is seen that the spin-dependent diffusion induces an effective depletion of majority electrons at the center due to an enhanced diffusion of these electrons. This spin filter effect produces a decrease of the corresponding spin polarization, in remarkable contrast with the situation under nondegenerate diffusion conditions in which the polarization will always decrease as a function of distance due to the spin relaxation in the semiconductor.

Shown in the top left panel of Fig. 6 are PL polarization profiles for selected excitation powers. These profiles reveal a transition into a degenerate regime which manifests itself by the progressive appearance of the dip at the center. In Curve f of Fig. 6 taken at the maximum power, the luminescence polarization
FIG. 5: 3-D representation (and its 2-D projection) of the electronic polarization for four different excitation powers at 15K. a) 72 nW, b) 0.45 mW, c) 1.89 mW, d) 2.55 mW. Qualitative explanation for the dip at the center in the polarization profile is given in the inset. This inset shows typical profiles of concentration of $n_{+}$ and $n_{-}$ electrons as a function of space, considering that $D_{++}$ is greater than $D_{--}$. Since majority electrons will diffuse further than minority carriers, a spin filter effect is predicted. A decrease of the electronic polarization (dashed line) at the excitation spot is expected, together with the appearance of a maximum at some distance away from the center, before the decrease to zero as a result of spin relaxation.

at $r = 0 \mu m$ is 13 %, smaller than its value of 19 % at $r \approx 2 \mu m$.

**B. Charge/spin coupling**

In order to reveal the second order effect predicted by Eq. (26), according to which the charge diffusion constant depends on spin, two sum profiles were compared. The first one $I_{\pi}$ was taken for spin-unpolarized photoelectrons generated by $\pi$ excitation. Keeping the excitation power constant to about 0.1%, a second sum profile sum profile $I_{\sigma}$ was taken for spin-polarized photoelectrons generated by $\sigma$ excitation. Figure 7 shows the relative difference of these profiles at $T_{e} = 15 K$ for different power densities. Curve d, shows that, as expected, charge diffusion does not depend on spin at low power, that is, in non degenerate conditions. In contrast, when the excitation power is increased, there progressively appears a spin dependent diffusion resulting in a depletion of photoelectrons at the center. Such depletion is compensated by a converse excess of photoelectrons at a distance larger than about 1.5 $\mu m$. This shows, in agreement with Eq. (26), that the spin-averaged diffusion constant is larger for spin-polarized electrons.

**C. Temperature dependence of the spin-dependent diffusion**

As seen from Fig. 1 and Eq. (5), since the effective density of states increases with temperature, the above effects are expected to be observable only at low temperature. Figure 8 shows the polarization profile at high excitation power (2.55 mW) taken at different sample temperatures, corresponding to different temperatures of the electron gas $T_{e}$. The dip at the center due to Pauli...
V. INTERPRETATION

A. Qualitative analysis at the center

It is shown here that all the experimental results are in excellent agreement with a qualitative analysis, considering only Pauli blockade effects and assuming for simplicity that charge and spin concentrations in steady-state are independent on $z$. This analysis considers values of luminescence intensity and polarization at the center, $r=0$, at which it can be shown using Sec. 11.13 that, in the same way as for Ref. 22, the spin-averaged lifetime of the electrons $\tau_{eff}^0$ is limited by diffusion rather than by bulk or surface recombinations. As a result, within numerical factors of order unity, one has

$$\tau_{eff}^0 \approx \omega^2/[4D_n\nu(n/2)] \sim 10^{-11}\text{s}$$

where $\omega$ is the radius of the Gaussian laser spot. The characteristic excitation power $P$ required for degeneracy at 15 K is then given by $g(0)\tau_{eff}^0 > N_e$. At the temperature of the electron gas of about 50 K which is measured, one has $N_e \approx 2 \times 10^{10}$ cm$^{-3}$. Thus, one finds $P \geq 1$ mW, in good agreement with the transition observed between curves $c$ and $d$ of Fig. 6. The measured value of the electron polarization at the center also gives an estimate of the ratio $D_{++}/D_{--}$. Since at $r=0$, spin-lattice relaxation at the bottom of the conduction band is negligible, one can write using Eq. 37

$$n_{\pm} = g_{\pm}(0)^*\omega^2/[4D_n\nu(n_{\pm})]$$

where $g_{\pm}(0)^*$ are the rates of creations of spin-polarized electrons at the bottom of the conduction band. Taking account of the slight spin-lattice relaxation during thermalization, one has $g_{\pm}(0)^*/g_{\pm}^*(0) = (1 + \mathcal{P}(0))/[1 - \mathcal{P}(0)]$, where $\mathcal{P}(0)$, of 0.45, is the polarization measured at the center at low excitation power. Eq. 38 gives the following very simple result

$$D_{++}/D_{--} = \frac{1 + \mathcal{P}(0)}{1 - \mathcal{P}(0)} \times \frac{1 - \mathcal{P}(0)}{1 + \mathcal{P}(0)}$$

At low power no dip is observed, so that $\mathcal{P}(0) = \mathcal{P}(0)$, and the ratio $D_{++}/D_{--}$ is indeed found equal to unity. At high power density, the latter equation gives $D_{++}/D_{--} \approx 1.54$. In order to estimate $D_{++}/D_0$ and $D_{--}/D_0$ separately, we numerically solve Eqs. (33) using the linear regime given by Eq. (11). This gives at $T_e = 50$ K $n_+ \approx 4.6 \times 10^{16}$ cm$^{-3}$ and $n_+ \approx 2.7 \times 10^{16}$ cm$^{-3}$. Comparison between these values and the spin-resolved effective density of states shows that, in fact, the observed very large spin dependence of the diffusion corresponds to the onset of degeneracy. Using Eqs. (12) and (33), one obtains $D_{++} = 3.4 D_0$ and $D_{--} = 2.3 D_0$. The mobility is also found to depend significantly on spin, since $\mu_+ = 1.7\mu_0$, $\mu_- = 1.4\mu_0$, so that $\mu_+ / \mu_- \approx 1.2$. 

blockade indeed decreases with increasing $T_e$, suggesting a transition from a non-denegerate to degenerate photoelectron gas between $T_e = 90$ K and $T_e = 110$ K, corresponding to a sample temperature of $T = 60 - 80$ K.

![Figure 7](image1.png)

FIG. 7: Left panel: relative difference between the sum $n = n_+ + n_-$ profile obtained under circularly polarized excitation (σ) and under linearly polarized excitation (π), for different power densities, a) 2.33 mW, b) 0.95 mW, c) 0.41 mW and d) 65 nW. For each excitation power, the only difference is the polarization of the created photoelectron gas. A difference of the order of 2.5 % between both profiles is observed at high power at $r=0$, revealing the spin-dependent diffusion of photoelectrons. Right panel: prediction given by a numerical solution of Eqs. (19) and (20) with no adjustable parameters.

![Figure 8](image2.png)

FIG. 8: Polarization profile at high power density ($P = 2.55$ mW) for different sample temperatures, corresponding to different values of the electronic temperature $T_e$, a) 55 K, b) 57 K, c) 68 K, d) 82 K, e) 90 K, f) 110 K, g) 200 K, h) 360 K. A clear transition from the Pauli blockade regime is observed with increasing $T_e$ between $T_e = 90$ K (curve e) and $T_e = 110$ K (curve f).
We now interpret the magnitude of the spin dependence of the charge diffusion constant [Fig. 7]. The spin-dependence of the charge concentration at the center is given by

\[
\frac{n_\sigma - n_\pi}{n_\sigma + n_\pi} = \frac{\langle D_\sigma \rangle - \langle D_\pi \rangle}{\langle D_\sigma \rangle + \langle D_\pi \rangle} = -\frac{\eta \mathcal{D}^2}{2 + \eta \mathcal{D}^2} \quad (40)
\]

For \( n \approx 7.3 \times 10^{10} \text{ cm}^{-3} \) calculated above, one finds neglecting spin-drag, \((n_\sigma - n_\pi)/(n_\sigma + n_\pi) \approx -2.7 \% \). This very simple estimate at the center gives a value in very good agreement with the experimental one.

We finally interpret the results as a function of temperature, summarized in Fig. 8. Using the above estimate of \( n \) at the center, (almost independent of temperature since \( D_0 \) weakly depends on \( T \)), the electronic temperature that determines the transition into a degenerate regime is found to be \( T_e = 300(g\tau_{eff}/N_0^0)^{2/3} = 110 \text{ K} \), where \( N_0^0 \) is the effective density of states density at room temperature. This is also in excellent agreement with the temperature dependence observed on Fig. 8.

B. Quantitative analysis using the coupled diffusion equations

For precise quantitative analysis, since we estimate \( \eta \mathcal{D} = 0.3 \), we have used the non linearized coupled Eqs. (10) and (17). These equations were solved self-consistently together with Eq. (9) using a commercial finite element package. The various parameters were not adjusted and were taken equal to their values given in Sec. III B. The spin diffusion constant was taken equal to their values given in Sec. III B. The spin diffusion constant was taken equal to the charge one. Integrations along \( z \) were performed in order to calculate experimentally measured profiles, as given by Eqs. (27), (28) and (29). In the bottom panel of Fig. 6 are shown the polarization profiles predicted by this model as a function of light power. As seen from a comparison between the bottom right panel, obtained by taking \( N_0^0 = 0.15N_0^0 \) in Eq. (22) and the bottom left panel corresponding to a unipolar model where one takes \( D_\pi = D_0 \) in the latter equation, the predicted magnitude of the effects is smaller in the latter case. Ambipolar diffusion does not affect the ratio \( D_{++}/D_{--} \), but results in a decrease of the effective global diffusion constant, and thus increases the number of photoelectrons at the center and Pauli blockade is enhanced. The numerical results show a maximum of polarization at \( r \approx 2 \mu \text{m} \), in excellent agreement with the experimental results. However, the peak value of the polarization predicted is even higher than the polarization at low power density at \( r = 2 \mu \text{m} \), whereas the experimental results reveal a global decrease of polarization as the power density is increased.

The latter difference is attributed to local heating of the electron gas due to the strongly-focused light excitation. Shown in Fig. 9 is the distance dependence of the electronic temperature. This figure reveals at very high power density a heating of the electron gas, for which the temperature is of the order of 90 K at the center and decrease to the value of 50 K, obtained for a weak excitation power, at a distance of about 2 \( \mu \text{m} \). Such heating can induce a decrease of the spin-lattice relaxation time \( T_1 \), and a polarization decrease. As seen in the inset of Fig. 9 such effect also manifests itself from the power-induced decrease of the spatially-averaged luminescence polarization \( \langle \mathcal{P} \rangle \). Complete inclusion of this effect would require to introduce a spatially-dependent value of \( T_1 \). Here, for simplicity, we consider a spatially-independent value of \( T_1 \). Since the spatially-averaged polarization comes preferentially from electrons near the center, it is reasonable to take the value which gives the measured value of \( \langle \mathcal{P} \rangle \) according to the usual relation \( \langle \mathcal{P} \rangle = \mathcal{P}_c^2T_1/(T_1 + \tau) \). Also shown in the inset of Fig. 9 are the values of \( T_1 \) used in the numerical calculation.

The results obtained this way, shown in the top right panel of Fig. 6 are in very good agreement with experimental results. The same method was used to calculate the spin-induced change of luminescence intensity and the results are shown in the right panel of Fig. 7 along with the experimental results shown in the left panel. Again, a very good agreement is obtained.
VI. DISCUSSION

A. Spin-Drag

Since spin-drag and Pauli blockade appear as additive contributions in Eq. (20) their effect on the polarization profile could be qualitatively similar. The calculated ratio of the off-diagonal to diagonal elements of the resistivity matrix, in conditions similar to the present experimental situation, is shown in Fig. 2 of Ref. [15]. This shows that \( \tau_m/\tau_{ee} \) is of the order of 0.2 - 0.3.

In the present subsection it is shown that spin-drag alone cannot explain the observed effects. Firstly, the spin dependence of the intensity profile (Fig. 7) cannot be due to spin drag, since spin drag does not appear in the charge diffusion equation [ Eq. (19)]. The results of Fig. 4 are thus a direct proof of Pauli blockade. The second evidence for the secondary role of spin-drag comes from calculations. Under the assumption that \( \tau_m/\tau_{ee} \) is constant in space and including only the effect of spin-drag, the spin diffusion equation becomes

\[
\left[ g \mathcal{P}_s - \frac{\tau_m}{\tau_m + \tau_{ee}} \left( g - \frac{n_m}{\tau_s} \right) \right] s + s D_0 \left( \frac{\tau_e}{\tau_m + \tau_{ee}} \right) \nabla^2 s = 0
\]

(41)

Shown in Fig. 10 is a numerical calculation of the spatial dependence of the polarization. The dotted lines represent the spatial polarization profile calculated from Eqs. (19) and (41). The parameters used in the calculation are again those obtained in Sec. III B and the four subfigures correspond to different values of the ratio \( \tau_m/\tau_{ee} \). From Eq. (41) it is seen that at sufficiently large distances, so that \( \nabla^2 s = 0 \) and \( g = 0 \), the spin density \( s \) is proportional to the charge density \( n_m \), so that the polarization becomes constant in space. Also, the first term of Eq. (41) can be viewed as an effective creation rate smaller than \( g \mathcal{P}_s \), which induces a drop in the polarization at the center. Both features are confirmed by the dotted lines of Fig. 10.

A more realistic model consists in solving Eqs. (19) and (20), considering only spin drag and assuming, in overall agreement with theoretical results [16], that \( \tau_m/\tau_{ee} \propto n_m \) and that using selected values of the ratio \( \tau_m/\tau_{ee} \) at \( r = 0 \). The results are shown by the solid lines of Fig. 10. The polarization profiles are found to be very different from those shown in Fig. 6. A polarization dip is observed at extremely high values of \( \tau_m/\tau_{ee} \geq 0.5 \), but its position and depth are very different from those of Fig. 6.

In summary, although it cannot be excluded that spin-drag can play a role in the polarization profiles of Fig. 6, it has been shown here that the dominant effect observed here for spin-dependent diffusion is indeed Pauli blockade.

![Fig. 10: Luminescence polarization profiles obtained by a numerical resolution of Eqs. (19) and (20) which takes into account only the spin-drag effect. The different values of \( \tau_m/\tau_{ee} \) used are a) 0.1, b) 0.5, c) 0.7 and d) 1. The dotted lines correspond to a model where this ratio is constant in space. The polarization decreases at the center and acquires a constant value at large distances, as predicted by Eq. (11). In contrast, if one assumes \( \tau_m/\tau_{ee} \propto n_m \), the solid lines are obtained, and a polarization dip can be observed at very high values of \( \tau_m/\tau_{ee} \). However, neither the position of the maximum, nor the values of the polarization are in agreement with the experimental profiles of Fig.](image-url)

B. Position of the polarization maximum

It is anticipated that the distance \( r_m \) from the center at which the polarization maximum occurs should be of the order of the effective spin diffusion length. However, analytical verification of the latter hypothesis is not simple due to the very large variation of the electron concentration between 0 and \( r_m \). Using Eqs. (33) and (21), we expect that \( r_m \) is approximately given by

\[
r_m = L_s \sqrt{\nu (n^*/2)}
\]

(42)

where \( n^* \) is a typical value of the electronic concentration, between \( r = 0 \) and \( r_m \). In order to verify the proportionality between \( r_m \) and \( L_s \), we have reproduced the numerical calculations of the preceding subsection, by varying the diffusion constant, which ensured a constant spatially-averaged polarization, and by adjusting the excitation power to ensure a constant electron...
concentration at the center. The heating of the electron gas was not included. Shown in the top left panel of Fig. 11 are polarization profiles, calculated using a concentration at the center of \( n \approx 7.3 \times 10^{16} \text{cm}^{-3} \). While Curve b reproduces the behavior of the bottom left panel of Fig. 6, the other curves were calculated using a diffusion constant \( D_0 \) multiplied by 0.1 (Curve a), 30 (Curve c), and 100 (Curve d). One sees that, as expected, \( r_m \) increases with \( D_0 \). Shown by open squares in the right panel of Fig. 11 is the dependence of \( r_m \) as a function of the spin diffusion length \( L_s \). The dependence is indeed linear. As shown from the dotted line, the value of \( r_m \) is smaller than \( L_s \) by about 20\%, implying that \( n^* \) is close to the concentration value at \( r_m \).

The bottom left panel of Fig. 11 was calculated for an electronic concentration larger by one order of magnitude. As expected, the relative height of the polarization maximum is strongly increased. Also visible is the fact that \( r_m \) is increased, as expected from Eq. \( (12) \). The dependence of \( r_m \) as a function of \( L_s \) is also shown in the right panel of Fig. 11. The quantity \( r_m \) still depends linearly on \( L_s \). The increase of \( r_m \) is about a factor 3, which is qualitatively similar to the ratio of the quantities \( n \) at the center for the two concentrations, of 4.8. It is concluded that the position of the polarization maximum, within numerical factors of order unity, corresponds to the effective spin diffusion length \( L_s \sqrt{n^*(n^*/2)} \).

VII. CONCLUSION

In the present work, we present an experimental investigation of the charge and polarization profiles of photoelectrons of a p+ GaAs film, using a microluminescence setup in which a circularly-polarized laser is tightly focused to a spot much smaller than charge and spin diffusion lengths. In conditions where the photoelectron gas is degenerate, i.e., for a sufficiently low temperature (near 15K) and large excitation power, polarization profiles exhibit a maximum at some distance away from the excitation spot, in strong contrast with the usual decrease of polarization caused by spin-lattice relaxation. This effect reveals that diffusive transport depends on spin, since such spin-dependent transport creates a spatial spin filter effect which results in an effective increase of polarization during diffusion.

Relative differences in the spin-resolved diffusion constants as large as 55\% are found between the two types of spins. This spin dependence is the product of the relative spin dependence of the mobility and of that of the proportionality coefficient between diffusion constant and mobility. Both these contributions are equivalent and of the order of 20\%. The spin-averaged charge diffusion constant is also shown to be spin-dependent, since in degenerate conditions, the charge spatial profile is shown to depend on the average spin polarization of the photoelectrons, due to a second order effect.

This novel spin-charge coupling mechanism is a direct manifestation of Pauli blockade in transport of a Fermion gas which forbids elementary scattering mechanisms. It is distinct from the previously-known spin-drag mechanism, which cannot explain the spin filter effect and the spin dependence of the charge diffusion constant. In contrast, a simple model including the Pauli blockade and ambipolar diffusion quantitatively accounts without any adjustable parameter for all the experimental results, including their dependence as a function of temperature and excitation power.

It is pointed out that these large spin-dependent effects have been observed in a regime near the onset of degeneracy, where the photoelectron concentrations are not very large with respect to the effective density of states in the conduction band. This implies that much stronger effects are expected for larger powers. While this is not possible in the present case because of heating effects, we anticipate that the use of appropriate low dimensional structures of reduced effective density of states will allow us to increase the magnitude of the effects and to possibly open the way to the realization of spin components of increased diffusion length and mobility at a temperature closer to 300 K.

![FIG. 11: The left panels show, for two electron concentrations at the center, the calculated polarization profiles for increasing values of the diffusion constant. These values are the actual values of \( D \) for curve b, while they are reduced by 10 for curves a, and multiplied by 30 for curves c and 100 for Curves d. It is seen that the position \( r_m \) of the polarization maximum shifts with increasing \( D \) as well as with increasing electron concentration. These results are summarized in the right panel which shows that \( r_m \) is approximately proportional to the spin diffusion length \( L_s \) given by Eq. \( (53) \), for the smaller (open squares) and larger (full circles) electronic concentration. As shown by open circles, the values at high concentration, correspond to \( r_m \) at low concentration multiplied by a factor of 3.](image-url)
Acknowledgments

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