Chapter 5

Spin Relaxation

5.1 Introduction

Spintronics [1,2] (or "spin-based electronics"), also known as magnetoelectronics, is an emergent technology which exploits the quantum spin states of electrons as well as making use of their charge state. It is a nontrivial extension of conventional electronics, augments functionality by utilizing the carrier spin degree of freedom. Spin can potentially be used as a much more capacious quantum information storage cell, be involved in the transfer of information, for elaborate schemes of information processing, both quantum mechanical and classical, and be integrated with electric charge counterparts in combined designes. The electron spin itself is manifested as a two state magnetic energy system. The discovery of Giant Magnetoresistance in 1988 by Albert Fert et al. [3] is considered as the birth of spintronics for which Nobel prize has been awarded in 2007.

In addition to their mass and electric charge, electrons have an intrinsic quantity of angular momentum called spin, almost as if they were tiny spinning balls. Associated with the spin is a magnetic field like that of a tiny bar magnet lined up with the spin axis. For a sphere spinning "west to east" the spin vector points "up" as shown in Figure 5.1. It points "down" for the opposite spin. In a magnetic field, electrons with "spin up" and "spin down" have different energies. In an ordinary electric circuit the spins are oriented at random and have no effect on current flow. Spintronic devices create spin-polarized currents

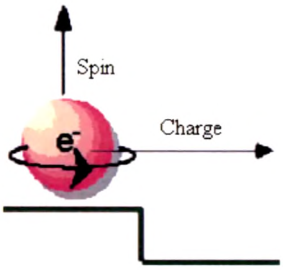

Fig. 5.1 Spinning Sphere of electron consisting spin and charge.

and use the spin to control current flow. Spintronic devices are used in the field of mass-storage devices and recently IBM scientists are successful in compressing the massive amounts of data into a small area, at approximately one trillion bits per square inch (1.5 Gbit/mm^2) or roughly 1 TB on a single sided 3.5" diameter disc. The ultimate goal in the rapidly burgeoning field of spintronics is to realize semiconductor based devices that utilize the spin degree of freedom of electrons instead of or in addition to charge degree of freedom [4]. Scientists are exploring the field of spintronics, in which the spin degree of freedom of the electron can be utilized to sense, store, process and transfer information. Taking advantage of the spin one can not only improve the performance of and add new functionality to existing devices but also could revolutionize electronics, leading to a new generation of multifunctional spintronic devices. There exist myriad spintronic device proposals, all of which rely on *coherent* spin transport in quantum confined semiconductor structures [5, 6]. An extremely popular representative of this huge class of (proposed) spintronic devices is the so-called "spin field effect transistor" or spin – FET [7]. In this device, electrons are injected into a semiconductor channel from a ferromagnetic (or half metallic) source contact with a definite spin polarization. The spin is then controllably precessed in the channel with a gate potential using the Rashba spin-orbit coupling effect [8]. Finally, when the electron arrives at the ferromagnetic/half metallic drain contact, it is transmitted with a probability that depends on the orientation of the electron spin with the fixed magnetization direction of the drain. By changing the probability with the gate potential that precesses the electron spin in the channel, one can modulate the channel current and realize transistor. However, there exist many major roadblocks for realizing a spin - FET or any other spintronic device. Particularly, electron spins, once injected into the semiconducting channel, experience in their own frame of reference a magnetic field which is generally referred as *spin-orbit coupling* [9-16]. This effect tends to orient the spins of different electrons in different directions and the spin coherence of the ensemble is lost. This phenomenon is commonly referred as "spin relaxation" (SR) and is detrimental to spintronic devices. Thus it is essential to have a proper understanding of spin-orbit mechanism-mediated spin relaxation processes that are operative in semiconductor nanostructures.

There are several mechanisms that can cause spin relaxation of conduction electrons:

(i) The D'yakonov-Perel' mechanism [17, 18] arises because spin-orbit interaction results in a wavevector dependent magnetic field **B(k)** about which the electron
(ii) spin precesses. It regards the spin splitting of the conduction band in zinc-blende semiconductors at finite wave vectors as equivalent to the presence of an effective magnetic field that causes electron spin to precess. For an electron experiencing random multiple scattering events, the orientation of this effective field is random, thus leading to spin relaxation.

(iii) On the other hand, Elliott-Yafet (E-Y) [19] mechanism occurs because in presence of spin-orbit interaction spin states are momentum dependent. Any spin-independent momentum scattering event has a finite probability to flip spin causing spin relaxation.

(iv) Bir-Aranov-Pikus Mechanism (BAP) [20] is the exchange interaction between electrons and holes which leads to spin relaxation in p-doped semiconductors. In this type of mechanism, simultaneous flip of electron and hole spins occur due to electron – hole exchange coupling.

(v) Rashba mechanism [21] is similar to the D-P mechanism in which an asymmetry is induced by due to the external electric field which causes the carriers to lose their spin orientation.

Spin relaxation can also take place as a result of hyperfine interaction of electron spins with magnetic momenta of lattice nuclei, the hyperfine magnetic field being randomly changed migration of electrons in the crystal. The thorough discussion of these mechanisms is in the following section.

5.2 Spin Relaxation Mechanisms

It was observed around 1967 by A. Kastler [22] and his followers that the application of the optical orientation method to semiconductors creates the unique opportunity to measure very short spin relaxation times of electrons and holes, down to 10^{-12} s, inaccessible to the traditional ESR techniques. This method allows studying spin relaxation of both free and bound carriers in crystals with different doping levels and in a wide temperature range. Consequently, very interesting new experimental data were available at that time which can't be explained in the frame work of spin relaxation mechanisms known at that time. This situation stimulated further theoretical studies which have brought above mentioned mechanisms.

The lake of an inversion centre in some semiconductor compounds leads to spin splitting of the conduction band. D'yakonov and Perel' [23] have proposed a relaxation mechanism due to this splitting. They have analysed this relaxation mechanism in terms of motional narrowing. In optical orientation experiments one studies SR of non-equilibrium carriers, usually photoelectrons in p-type crystals. In the case of high hole concentration another relaxation mechanism, due to exchange interaction between electrons and holes, can be very efficient which is known as Bir-Aranov-Pikus Mechanism (BAP); detailed description of which has been given by Bir et al. [24]. In some cases the Elliot- Yafet spin relaxation mechanism (Elliot 1954, Yafet 1963) can also be important in optical orientation experiments. In this mechanism spin relaxation of carriers is connected with momentum relaxation through spin - orbit coupling. The Rashba effect [25], named after one of its discoverers, is an energy splitting, of what would otherwise be degenerate quantum states, caused by a spin-orbit interaction in conjunction with a structural-inversion asymmetry in the presence of interfacial electric fields in a semiconductor heterostructure. Under certain conditions in semiconductors an appreciable reabsorption of recombination radiation with simultaneous creation of new electron-hole pairs can take place. This reabsorption process causes a decrease of photoelectron spin orientation and can be treated also as a spin relaxation process.

In the following we shall present the theory of above mentioned spin relaxation mechanisms. Special attention will be given to spin relaxation under uniaxial compressive stress and in a longitudinal magnetic field with the restriction to one group of crystals.

5.2.1 The Elliott-Yafet Mechanism (EY Mechanism)

As was pointed out by Elliott (1954) and later by Yafet (1963), spin-orbit interaction in crystals causes the mixing of electron wave functions with opposite spin vector orientations. As a result, in the process of momentum scattering the disorientation of electron spin also becomes possible. Afterwards, the spin relaxation rate for this mechanism has been calculated in detail by number of authors (Pavlov and Firsov 1965, Pavlov 1966, Abakumov and Yassicvich 1971, Chazalviel 1975). In these calculations different momentum relaxation processes have been taken in account. In the following subsections we shall, in short, present calculations of the spin relaxation rate for the cases of electron scattering by acoustic and optical phonons and by impurities.

5.2.1 (a) Long - Range Interaction

Considering the mixing of wave functions of the Γ_7 band with wave functions of the Γ_8 and Γ_6 bands, we can note the spin flip matrix element for functions between m and m' states in the conduction band as follows:

$$\mathbf{H}_{m'k',mk} = \sum_{s,s'} \frac{\mathbf{H}^{k}_{m'k',s'k'} \mathbf{H}^{\text{int}}_{s'k',sk} \mathbf{H}^{k}_{sk,mk}}{\left(E_s - E_m\right)\left(E_{s'} - E_m\right)} \tag{5.1}$$

Here, $\mathbf{H}^{\mathrm{k}}_{\mathrm{sm}} = \mathbf{H}^{k'}_{ms}$ are the terms linear in k. and are responsible for the mixing of wave functions, and $\mathbf{H}^{\text{int}}_{s',s}$ are the matrix elements of the operator which determines the electron interaction with phonons or impurities. The indices m and s refer to states in conduction and valance bands, respectively. The Polar long-range interaction with optical phonons is described by Fröhlich Hamiltonian

$$D^{opt}_{k',k} = iC\frac{q \cdot U}{q^2} I\delta_{q,k-k'} \tag{5.2}$$

where I is the unit matrix, $C = e\omega_{LO} \times \left(4\pi\rho_0 / \varepsilon^*\right)^{1/2}$, $\varepsilon^{*-1} = \varepsilon_{\infty}^{-1} - \varepsilon_{0}^{-1}$ where ε_0 and ε_∞ are the static and high frequency dielectric constants respectively, ρ_0 is the crystal density, ω_{LO} and $\mathbf{q}$ are the frequency and wave vector of the longitudinal phonon, $\mathbf{k}$ and $\mathbf{k}'$ are the electron wave vectors for initial and final states respectively. The choice of the constant C in above equation corresponds to the following definition of U:

$$U=\frac{2\sqrt{M_1M_2}}{M_1+M_2}(U_1-U_2) \tag{5.3}$$

where U_1 and U_2 are the displacements of the sublattices formed by atoms with masses M_1 and M_2 respectively. Long range electron-phonon piezoelectric interaction in crystals with T_d symmetry is described by Hamiltonian

$$\mathbf{D}^{\mathrm{ac}}_{\mathrm{k',k}}=i\frac{8\pi e\beta_0}{\varepsilon_0 q^2}\left(q_x\varepsilon_{yz}+q_y\varepsilon_{zx}+q_z\varepsilon_{xy}\right)I\delta_{q,k-k'}, \tag{5.4}$$

where β_0 is the piezoelectric modulus.

Interaction with ionized impurities is described by Hamiltonian

$$V_{k',k}=\frac{4\pi e^2}{\varepsilon_0}\frac{1}{q^2+L_D^{-2}}I\delta_{q,k-k'} \tag{5.5}$$

where L_D is the Debye screening length.

According to eqs. (5.2)- (5.5) for long range interaction,

$$\mathbf{H}^{\mathrm{int}}_{s',k'}=\mathbf{H}^{\mathrm{int}}_{k',k}\cdot\delta_{s,s'} \tag{5.6}$$

Now, the

$$\mathbf{H}_{k',k}=\mathbf{H}^{\mathrm{int}}_{k',k}\frac{i\hbar^2\left(\sigma[k'k]\right)\eta\left(1-\frac{1}{2}\eta\right)}{3m_cE_g\left(1-\frac{1}{3}\eta\right)} \tag{5.7}$$

For elastic scattering, when $\mathbf{k'}=\mathbf{k}$,

$$\left(\frac{\hbar^2}{m_e}[k'k]\right)^2=4E_k^2\sin^2\theta$$

where, $E_k = \hbar^2 \mathbf{k}^2 / 2m_e$ is the electron kinetic energy, θ is the scattering angle. The transition probability should be averaged over all directions of **k** as the spin relaxation rate of electrons isotropically distributed in k space. The spin relaxation rate is given by

$$\frac{1}{\tau_s} = 2\frac{2\pi}{\hbar}\int\frac{d\Omega_k}{4\pi}\sum_{\mathbf{k}'}\left|H_{m'\mathbf{k}',m\mathbf{k}}\right|^2 \delta\left(E_{\mathbf{k}'}-E_{\mathbf{k}}\right) \quad ; \quad \left(m' \neq m\right) \qquad (5.8)$$

The simplification of above eq. (5.8) leads to the following equation of the spin relaxation rate,

$$\frac{1}{\tau_s} = \frac{16}{27}\frac{1}{\tau_p}\left(\frac{E_k}{E_g}\right)^2 \eta^2 \left(\frac{1-\frac{1}{2}\eta}{1-\frac{1}{3}\eta}\right)\Phi \qquad (5.9)$$

where,
$$\Phi = \frac{\int_{-1}^{1}\left(1-\mu^2\right)\sigma(\mu)d\mu}{\int_{-1}^{1}\left(1-\mu\right)\sigma(\mu)d\mu}$$

Here $\mu = \cos\theta$ and σ(μ) is the scattering cross section. For scattering by polar optical phonons, if the electron kinetic energy E_k is much higher than the LO phonon energy $\hbar\omega_{LO}$, we can obtain σ(μ)=(1-μ) and Φ=1. For scattering by ionized impurities $\sigma(\mu) \approx \left[(1-\mu)+2/\beta\right]^{-2}$, where $\beta = \left(2kL_D\right)^2$ and

$$\Phi = \frac{2(2+\beta)}{\beta}\frac{\ln(1+\beta)-2\beta/(2+\beta)}{\ln(1+\beta)-\beta/(1+\beta)}$$

After averaging over the Maxwell distribution of the electrons the general expression for the spin relaxation rate due to EY mechanism can be written as

$$\frac{1}{\tau_s} = r\frac{1}{\tau_p}\left(\frac{k_B T}{E_g}\right)^2 \eta^2 \left(\frac{1-\frac{1}{2}\eta}{1-\frac{1}{3}\eta}\right) \tag{5.10}$$

where,

$$r = \frac{16}{27}\Phi\frac{\left\langle E_k^2 \tau_p^{-1}\right\rangle\left\langle E_k \bar{\tau}_p\right\rangle}{\left\langle E_k\right\rangle\left(k_B T\right)^2} = \frac{32}{81}\Phi\frac{\left\langle E_k^2 \tau_p^{-1}\right\rangle\left\langle E_k \bar{\tau}_p\right\rangle}{\left(k_B T\right)^3},$$

where, the angle brackets denote the averaging over the Maxwell distribution. Hereafter it is considered that τ_p , which determines the electron mobility $\mu_e = {e\tau_p}/{m_e}$, is given by

$$\left\langle \bar{\tau}_p E_k \right\rangle / \left\langle E_k \right\rangle = 2\left\langle \bar{\tau}_p E_k \right\rangle / 3k_B T$$

In the case of polar optical phonon scattering, r equals 2.0 and for the ionized impurity scattering $r = 3.0\Phi(\beta_1)$ where $\beta_1 = 4k_B T L_D m_e / \hbar^2$. For piezoelectric scattering, the accurate evaluation of $1/\tau_s$ and $1/\tau_p$ is rather difficult due to anisotropy of $D^{ac}_{k',k}$ in eq. (5.4).

5.2.1 (b) Short - Range Interaction

In the case of short range interaction for scattering rate by optical phonons it is sufficient to consider the admixture of wave functions of the Γ_8 and Γ_6 bands to only one of the wave functions of the Γ_7 band, since for this interaction scattering of electrons is possible from the Γ_8 or Γ_6 band to Γ_7 band. Therefore, for the interband scattering the matrix element between m',k' and m, k states are given as [26]

$$H_{m'k',mk} = -\sum_s \frac{\mathbf{H}^k_{m'k',sk'}\mathbf{H}^{\text{int}}_{sk',mk} + \mathbf{H}^{\text{int}}_{m'k',sk,}\mathbf{H}^k_{sk,mk}}{\left(E_s - E_m\right)} \tag{5.11}$$

Correspondingly, for the operator $H_{k',k}$, which gives the probability of spin- flip scattering, as given as ,

$$\mathbf{H}_{k',k} = -\frac{1}{3}\frac{\eta\hbar d_2\left(\sigma\left[U\times\left(k+k'\right)\right]\right)}{2m_e E_g\left(1-\frac{1}{3}\eta\right)^{1/2}} \tag{5.12}$$

In contrast to the case of long range interaction, scattering due to the short range potential occurs not only on longitudinal but also on transverse phonons. For the scattering on polar optical phonons, neglecting the difference in energies of longitudinal and transverse phonons, the following relation between $1/\tau_s$ and $1/\tau_p$ are given as [27]

$$\frac{1}{\bar{\tau}_s} = \frac{16}{27}\frac{1}{\bar{\tau}_p}\frac{E_k^2}{E_g E_0}\frac{\eta^2}{1-\frac{1}{3}\eta} \tag{5.13}$$

where $$E_0 = C^2\hbar^2 / 4d_2 m_e$$

In fact, above eq. (5.13) differs from eq. (5.9) only by the replacement of E_g^2 by $E_g E_0$. So the contribution of short range interaction exceeds that of long rage interaction when $E_0 > E_g$. The estimates shows that, for example, for GaAs this inequality holds if d_2a_0>40eV, a_0 being the lattice constant. Thus, short range interaction with optical phonons does not alter the temperature dependence of $\bar{\tau}_s$. A similar situation takes place fro the interband deformation potential scattering by acoustic phonons when the main contribution to $1/\tau_p$ is given by piezoelectric interaction. However, if the value of $1/\tau_p$ is determined also by the deformation potential interaction and interband contribution still prevails. The inverse spin relaxation time is given as [28]

Hence,
$$\frac{1}{\tau_s} = 0.088\frac{1}{\tau_p}\left(\frac{C_2}{C_1}\right)^2 \eta^2 \frac{k_B T}{E_g}\frac{1}{\left(1-\frac{1}{3}\eta\right)} \tag{5.14}$$

when the intraband contribution is dominating, calculation according to eq. (5.15) leads to following expression for $H_{k,k'}$

$$H_{k,k'} = i\left\{\left[\left(2-\frac{1}{3}\eta\right)a\varepsilon\left(\sigma[k'k]\right) - b\sum_i \sigma_i[k'k]_i\left(\varepsilon_{ij}-\frac{1}{3}\varepsilon\right) - \frac{d}{\sqrt{3}}\sum_{i\neq j}\sigma_i\varepsilon_{ij}[kk']_j\right\} \tag{5.15}$$

where a, b and d are the deformation potential constants for the valance band. Calculation according to eq. (5.8) with the substitution of eq.(5.15) for $\mathbf{H}_{m'k',mk}$ leads to a relation for $1/\tau_s$ similar to eq. (5.13) , with

$$\Phi = \frac{2\,\bar{a}^2}{3\,C_l^2} \tag{5.16}$$

where, $\bar{a}^2 = a^2 + \frac{1}{(2-\eta)^2}\left(\frac{8}{5}b^2 + \frac{3}{2}d^2 - \frac{1}{5}\sqrt{3}bd\right) + \frac{2a}{5(2-\eta)}\left(2b+\sqrt{3}d\right)$

In this case, $1/\tau_s$ is given by eq. (5.2) with $r = 2\left(\bar{a}/C_l\right)^2$.

The constant C_2 can be determined from the measurement of spin relaxation times in uniaxially deformed crystals. For GaAs it is approximately 3 eV. The smallness of C_2 in comparison with $\bar{a}$ is explained by the fact that the only reason for a non zero value of C_2 is the absence of inversion symmetry in the crystal.

5.2.2 The Bir-Aranov-Pikus Mechanism (BAP)

Electron scattering on holes can lead to spin-flip transition due to exchange and annihilation interaction. For electron-hole scattering, Bir, Aranov and Pikus in 1975, showed that the EY mechanism is less efficient than the BAP mechanism caused by the exchange interaction. Exchange interaction may also dominate in the case of electron scattering on paramagnetic impurities. In the following we shall briefly discuss the rate of electron spin relaxation due to BAP mechanism. However, there are many ways of scattering of electrons from holes such as free, bound etc., but in the present case, we discuss only one i.e. the scattering of electrons by free holes and this is the most important as well as mostly used for the calculation of spin relaxation rates [29]. The Hamiltonian for electron- hole exchange interaction is given by

$$H_{ex} = \pi a_B^3 D \delta(r) \delta_{k,k'} \qquad (5.17)$$

where $r = r_e - r_h$ is the position of the electron relative to the hole, $\hbar k$ and $\hbar k'$ are the total quasi momenta of the interacting electron and hole in the initial and final states, respectively. The factor, πa_B^3 where a_B is the exciton Bohr radius, is expressed in eq. (5.10) to express the operator D in energy units. In accordance with eq. (5.10), the electron spin relaxation time due to scattering on holes in the Bohr approximation is defined by

$$\frac{1}{\tau_s(k)} = 4\pi^3 a_B^3 \sum_{\substack{\alpha\alpha' \\ pq}} |\psi(0)|^4 \left| D_{m'\alpha', m\alpha}(p+q, p) \right|^2 f_p \left(1 - f_{p+q}\right) \times \delta\left(E_k + \varepsilon_p - E_{k-q} - \varepsilon_{p+q}\right)$$

$$(m' \neq m) \qquad (5.18)$$

Here, m and m' are the spin indices of an electron and α, α' are those of a hole in the initial and final states respectively, p and ε_p are the hole wave vector and energy $f_p(\varepsilon_p)$ is the hole distribution function, $\hbar q$ is the momentum transferred in the course of scattering. $\Psi(r)$ is the wave function of electron hole relative motion, $r=r_e-r_h$ is the position of the electron relative to the hole.

Let us now consider the most characteristic cases:

5.2.2 (A) *Holes are non-degenerate*

In this case at $m_e << m_h$ the transferred energy is smaller than E_k. Then, considering $\sum\limits_p f_p = N_p$ N_p being the hole concentration, eq. (5.11) for the spin relaxation rate can be reduced to the form,

$$\frac{1}{\tau_s} = \frac{2}{\tau_0}\frac{v_k}{v_B} N_p a_B^3 \qquad (5.19)$$

where,

$$\frac{1}{\tau_0} = \frac{\pi}{2\hbar}\frac{\overline{D}_s^2}{E_B}\left|\psi(0)\right|^4,$$

$$\overline{D}_s^2 = \int \frac{d\Omega_p}{4\pi} D_s^2(p) \qquad D_s^2(p) = \frac{1}{2}\sum_{qq'}\left|D_{m'q',mq}\right|^2 \qquad m' \neq m$$

$$E_B = \frac{\hbar^2}{2m_e a_B^2} \qquad v_B = \frac{\hbar}{m_e a_B} \text{ and } \qquad a_B = \frac{\hbar^2 \varepsilon_0}{c^2 m_e} \qquad (5.20)$$

5.2.2 (B) *Holes are Degenerate with Fermi energy*

In the case when holes are degenerate with Fermi energy $\varepsilon_F >> E_k, k_B T$ the spin relaxation rate can be given as [30]

$$\frac{1}{\tau_s} = \frac{3}{2}\frac{1}{\tau_0}\frac{\hbar^3}{m_c^2 v_B}\frac{N_p a_B^3}{\varepsilon_F}\int\frac{d^3q}{2\pi}\int_{-\infty}^{+\infty}\frac{\hbar\omega d\hbar\omega}{1-\exp\left(-\hbar\omega/k_B T\right)}$$

$$\times\int\frac{d\Omega_p}{4\pi}\delta\left(\hbar\omega-\hbar q v_F\cos\theta+\hbar^2q^2/2m_h\right)\delta\left(\hbar\omega-E_k+E_{k+q}\right) \tag{5.21}$$

The main contribution to $1/\tau_s$ is determined by scattering on heavily holes, since their concentration considerably exceeds that of light holes. Therefore, in the following we shall understand by m_h the mass of heavy holes and v_F their Fermi velocity. In this case

$$N_p=p_F^3/3\pi^2=\left((2m_h v_F)^{3/2}/3\pi^2\hbar^3\right)$$

For $E_k > m_e/m_h$, the electron energy E_e being smaller than ε_F. This inequality is satisfied at hole concentrations just above the critical concentration range can be called "The range of fast electrons", since under this condition the electron velocity exceeds that of holes at the Fermi surface. In this range $\hbar\omega=\hbar q v_F=\hbar k v_F\cong E_k v_F/v_k<<E_k$ as at $E_k\geq k_BT$ the transferred momentum $\hbar q$ does not exceed 2 $\hbar k$. At the same time $\frac{\hbar^2q^2}{2m_e}\frac{1}{\hbar\omega}=\frac{k}{p}<<1$. For thermalised and not very hot electrons, with $E_k<<(m_h/m_e)(k_BT)^2/\varepsilon_F$.

$$\hbar\omega=E_k\frac{v_F}{v_k}=\left(\varepsilon_F E_k\frac{m_e}{m_h}\right)^{1/2}<<k_BT$$

Here, $\hbar\omega/\left(1-\exp(-\hbar\omega/kT)\right)\approx k_BT$. After integration over q by integration over $k'\approx k-q$, neglecting $\hbar\omega$ in the argument of the last δ-function, the "fast" electron spin relaxation rate is given as

$$\frac{1}{\tau_s}=\frac{3}{\tau_0}\frac{v_A}{v_B}\frac{k_BT}{\varepsilon_F}N_p a_B^3 \tag{5.22}$$

For hot electrons, with $E_k >> (m_h/m_e)(k_B T)^2/\varepsilon_F$, the transferred energy satisfies the relation $\hbar\omega = E_k v_F/v_k >> k_B T$ and the hot electron spin relaxation rate is given as [30]

$$\frac{1}{\tau_s} = \frac{2}{\tau_0} \frac{v_F}{v_B} \frac{E_k}{\varepsilon_F} N_p a_B^3 \tag{5.23}$$

5.2.3 D'yakonov-Perel Mechanism

It seems to be well established that, among the several mechanisms leading to spin relaxation in GaAs quantum wells associated to the spin-orbit coupling in the electron dynamics, the dominant is the one known as the D'yakonov-Perel (D-P) mechanism. This mechanism appears in semiconductor lacking an inversion centre, as in the case of GaAs, where the splitting of the spin-polarized band is proportional to the cubic power of the quasi- momentum.

In the case of a quantum well the spin relaxation time (SRT), τ_s, given by the D-P mechanism is proportional to the inverse of the momentum relaxation time (MRT), τ_p. At low temperature this relation results in the fact that the cleaner the sample, the stronger the spin relaxation. The expression for the inverse of the SRT is:

$$\frac{1}{\tau} = \frac{\alpha^2 \hbar^2 \langle q^2 \rangle^2 k_B T}{2m^2 E_g} \tau_p \tag{5.24}$$

Where α is the coefficient for the spin splitting of the conduction band, m is the electron effective mass, E_g is the energy gap, k_B is Boltzmann constant, T is temperature and $<q^2>$ is the thermodynamic average of the square of the momentum in the confinement direction.

It has been observed experimentally [31] that the SRT in GaAs quantum wells is practically independent of the temperature in case of strong confinement (thin quantum wells). As we can see from eq. (5.24), besides the linear explicit dependence, the temperature appears implicitly in the MRT, as well as in the average of the kinetic energy in the direction of the confinement. We can take the MRT due to the scattering by impurities to be temperature independent (in case of non- magnetic impurities, at least). In consequence, the resulting behaviour with temperature coming out exclusively from that scattering mechanism is due to the linear dependence modulating the thermodynamic average of the perpendicular energy.

The work is motivated by the fact that, in general, the quasi-two- dimensional electron gas (2DEG) in GaAs quantum wells has its origin in remote doping. The electrons are, therefore, spatially separated from the ionized impurities generating them, in such a way that, in the temperature regime where the scattering by acoustic phonons is important, this mechanism dominates largely the scattering by impurities in the MRT. It is clear here that we are assuming the inverse of the observed MRT to be the sum of the inverses of the MRT due to every scattering source, the so- called Mathiessen's rule.

The MRT due to the scattering of the 2DEG by acoustic phonons via deformation potential can be easily obtained from the well known bulk results. In the quantum limit, where only the first subband is occupied at T=0, we have:

$$\frac{1}{\tau_p} = \frac{\varepsilon_{ac}{}^2 m}{\hbar^3 c_l L} k_B T \qquad (5.25)$$

In the above equation, L is the Quantum well width, ε_{ac} is the deformation potential, and c_l is the longitudinal elastic constant. This is similar to eq. (3.27) of momentum relaxation rate with D_A as acoustical deformation potential and W is well width.

We see that, if we neglect every other source of momentum relaxation, the linear modulation of the temperature cancels out the linear dependence of τ_p due to the scattering by phonons. The temperature dependence is totally content in the thermodynamic average of the perpendicular energy. The inverse of the SRT becomes:

$$\frac{1}{\tau_s} = \frac{2\alpha^2 \hbar L c_l}{m E_g \varepsilon_{ac}^2} \langle E_z \rangle^2 \tag{5.26}$$

Where the average perpendicular energy given by $< E_z >= {}_n E_n < N_n >$, where, $<N_n>$ representing the average number of electrons occupying the subband n, and E_n is the energy of the bottom of this subband. Assuming the confinement to be due to a very high barrier, $E_n = E_l n^2$, where, $n = 1, 2,3.....$, and $E_1 = \hbar^2 \pi^2 L^2 / 2m$. The average occupation number is given by

$$\langle N_n \rangle = \int_{E_n}^{\infty} d\varepsilon D_n(\varepsilon) f(\varepsilon, T), \tag{5.27}$$

With the subband density of states:

$$D_n(\varepsilon) = \frac{m}{\pi \hbar^2} \Theta(\varepsilon - E_n), \tag{5.28}$$

And the Fermi-Dirac distribution:

$$f(\varepsilon,T)=\frac{1}{\exp\left[\left(\varepsilon-\mu(T)\right)/k_B T\right]+1} \tag{5.29}$$

Using the constraint

$$n_s=\sum_n \int_{E_n}^{\infty} d\varepsilon D_n(\varepsilon) f(\varepsilon,T), \tag{5.30}$$

Where n_s is the areal density of the 2DEG, the chemical potential μ (T) can be easily obtained by numerically solving the equation:

$$\ln\left[\prod_{n=1}^{\infty}\left(1+\exp\left[\left(\left(\mu(T)-E_n\right)/k_B T\right)\right]\right)\right]=\frac{m}{\pi\hbar^2}n_s. \tag{5.31}$$

5.3 Results and Discussion

Our system consists of a GaAs quantum well confined by infinite barriers. In that case, α=0.028, ε_{ac}= 1.01×10^{-18} J, and c_f= 6.94×10^{10} N/m^2. We have considered, in our calculations, an electron gas with n_s=5×10^{11}cm^{-2}, and widths L=60,100,150,200 and 300 Å. The temperature from 50 to 300 K. Fig. 5.2 shows the renormalized inverse of the SRT as a function of the temperature. The renormalization is obtained by dividing the results of eq. (5.3) by the value obtained for L=100 Å at T=100K. The experimental results are plotted in the same graph, dividing energy value by the inverse of the SRT obtained in the sample of L = 10 nm at the temperature T= 87K. Our choice is made on the basis that these data seems to be more consistent with an SRT independent of temperature. Probably, the larger dispersion observed in the data corresponding to the sample with L= 6nm are due to reasons related to the preparation, which masks our conclusions. Of course, since we have neglected residual resistivity, the absolute value of SRT is probably meaningless, except for the fact that the residual resistivity does not change with temperature. On the other hand, the renormalized SRT reflects the real role of the temperature. Notice that the SRT is temperature independent for the width parameters L= 60 and 100Å, and changes for higher temperature for L= 150,200 and especially 300 Å.

It is worthwhile to know, for different temperature regimes, how the SRT changes with the quantum well width. This is shown in Fig. 5.3. The calculations are made for *T*=100K (solid line), 200K (dashed line), and *T*=300K (dotted line). The dependence with the quantum well width is very similar for this temperature. Then we plotted in the same graph the experimental values obtained in the four samples at *T*= 92 -100K (square), 200 K (circle) and 300 K (triangle). We observe a good agreement above 150 Å. For those samples corresponding to thinner quantum wells, the proximity of defects at the interface makes the scattering by these sources important enough not to negligible. The inverse of the MRT due to disorder becomes comparable to the acoustic phonons contribution. Since the total MRT is smaller than the one we calculate, discrepancies with our model are expected in very thin quantum wells. The inverse of the SRT we calculate must be higher than the one really observed. This is exactly what appears in Fig. 2. It must be said, however, that the scattering due to disorder dose not introduces temperature dependence on the SRT. To conclude, we have shown with a simple calculation that scattering by acoustic phonons via deformation potential leads to the experimentally observed behaviour of the SRT for electrons in quantum wells , in which reasonable thin widths result in temperature independent measurements , while for thicker wells an untypical temperature dependence shows up , as a transition to a three -dimensional behavior.

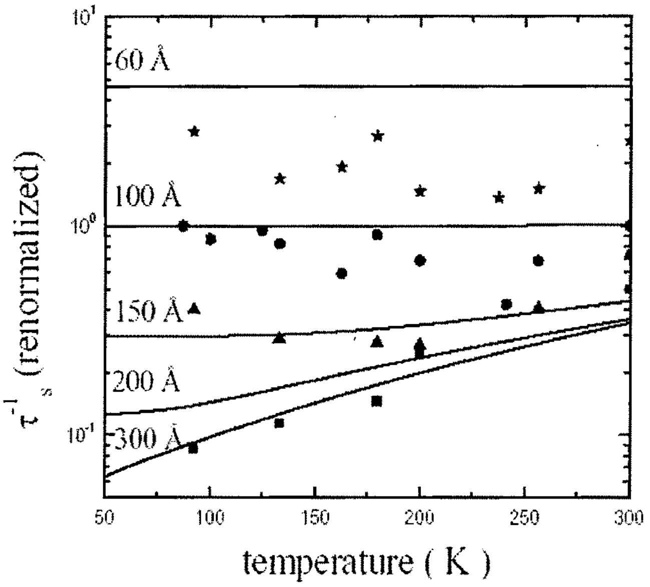

Fig. 5.2 The renormalized inverse of the SRT as a function of the temperature.

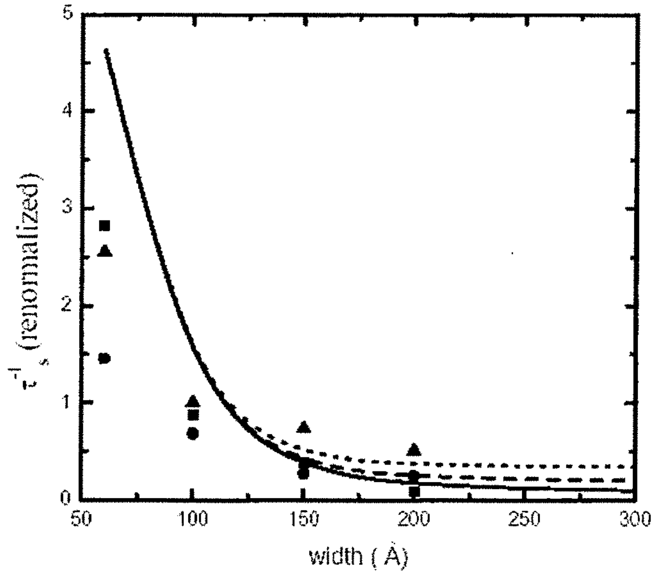

Fig. 5.3 The SRT changes with the quantum well width for different temperature regimes.

5.4 Summary

In this chapter, we have calculated the spin relaxation rate and by doing this we put one step ahead to understand the procedure of spin flip and hence improve the technology toward the mechanism of more little and fast memory storing devices. The chapter describes the different mechanisms of the spin relaxation. In the present formulation we use D'yakonov-Perel (DP) mechanism arising due to the spin-orbit interaction and are best suitable for the centro symmetry semiconductors lacking an inversion centre. We have shown that the scattering by acoustic phonons via deformation potential leads to the experimentally observed behavior of the **Spin Relaxation Time** for electrons in quantum wells, in which reasonable thin widths result in temperature independent measurements , while for thicker wells an untypical temperature dependence shows up , as a transition to a three – dimensional behavior. It is also observed here that the scattering due to disorder dose not introduces temperature dependence on the **Spin Relaxation Time**.

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