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TOPICAL REVIEW

The spin-valve transistor: a review and outlook

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Abstract

Combining ferromagnetic and semiconductor materials is a challenging route to create new options for electronic devices in which the spin of the electron is employed. The spin-valve transistor (SVT) is the first of such hybrid devices shown to work successfully. This review describes the basic science and technology of the SVT and derived devices, such as the magnetic tunnel transistor.

It seems almost inevitable that electronics in the near future will employ the electron spin degree of freedom. Two exciting scientific breakthroughs within the last 15 years and their exceptionally rapid implementation into products already demonstrates the impact of using the electron spin. In 1988, the giant magnetoresistance (GMR) effect [1] was discovered in multilayer structures that contain layers of ferromagnetic metals separated by a thin spacer of normal metal. The resistance of such structures was found to depend greatly on the relative magnetic orientation of neighbouring magnetic layers, making it attractive for application in highly sensitive magnetic field sensors [2–4]. As early as 1997, GMR was incorporated into the read-heads of magnetic hard disc recording systems, where it has been one of the main factors enabling the tremendous increase in storage density over the past decade. The second important breakthrough [5] was the demonstration in 1995 of large tunnel magnetoresistance (TMR) in tunnel junctions in which two ferromagnetic electrodes are separated by an ultrathin insulator. Key features [6] are the large spin dependence of the tunnelling process, even at room temperature, and the reproducible fabrication of reliable tunnel barriers typically using 1–2 nm of Al₂O₃. These features facilitated the development of a high performance magnetic random access memory (MRAM). It is non-volatile, has low power consumption and fast switching speed [7–9], and has the potential of becoming a universal memory. The first introduction of tunnelling based MRAM into the market is imminent.

Spin-transport in (quantum-) structures built exclusively from non-magnetic or paramagnetic semiconductors is explored actively, but a concern is that the creation of a spin polarization at present involves optical techniques or large magnetic fields to induce a sizable Zeeman splitting of the electron spin. It is thus almost inevitable that electronics in the near future will employ the electron spin degree of freedom.
spin states. It would be advantageous to use conventional magnetic materials such as Fe, Ni, Co and their alloys, in which ferromagnetism is robust and Curie temperatures are high. This supports a stable, virtually ever lasting memory function, while switching between different magnetization states is extremely rapid, typically at nanosecond or even picosecond timescales. It is, thus, attractive to examine different ways (figure 1) to combine ferromagnets and semiconductors.

The most straightforward approach is the one employed in MRAM, an array of magnetic memory elements is placed on top of a semiconductor wafer containing transistors and other circuitry required to drive the memory. This requires a modest, yet not trivial integration at the system level, but does not take advantage of the unique properties of semiconductors in manipulating spin. The most intimate form of integration is to put magnetic properties into semiconductor materials, thus creating ferromagnetic semiconductors [12]. Such materials can be obtained by doping with a certain amount of magnetic atoms, as in the case of the notable example of GaMnAs. Interplay between ferromagnetism and carrier densities via electrical gating has been achieved in some of the materials [13,14]. The hunt is now on for compounds that exhibit both semiconducting and ferromagnetic properties at temperatures well above room temperature.

An intermediate option is to create hybrid device structures in which ferromagnets and semiconductors are combined. One could think of many different geometries of such hybrid structures. Two main categories will be distinguished, based on whether the control and manipulation of the spins occurs in the semiconductor or in the ferromagnetic material. In the first category, electron spins that originate from a ferromagnetic source material are injected into a semiconductor, in which they are transported and manipulated, followed by some means of spin detection at the ‘other end’ of the device [15]. Much progress has recently been made on the first important step of spin-injection into the semiconductor [16–20], but implementation into working devices that operate at room temperature remains to be demonstrated. The opposite is true for the second class of hybrid devices, where a device concept has been successfully demonstrated by Monsma et al [21] with the introduction of the spin-valve transistor (SVT) in 1995, and the subsequent observation of huge magnetic response at room temperature [22,23] a few years later. While the spin dependence of the transport is in the ferromagnetic materials, the semiconductors are used to create energy barriers in the electron’s potential landscape that are essential to the operation of the device.

This review will cover this second category of semiconductor/ferromagnet hybrid structures, including the SVT and derived hot-electron devices such as the magnetic tunnel transistor (MTT). Hopefully, this review presents a basic understanding of such devices and provides insight into their strengths and weaknesses, which is also relevant for spin-electronics in general. Other spintronic devices, such as those based on spin-injection into semiconductors, in the end also have to address issues such as signal levels, noise, room temperature operation, device scaling, etc. It is important to realize the magnitude of the challenge this poses, as contemporary semiconductor electronics is a highly mature technology with exceptional performance. Spin-electronics will be successful only if it results in greatly enhanced performance, or, in a completely novel electronic functionality. Hence, merely creating a transistor based on spin (or carbon nanotubes or molecules) is not sufficient as silicon transistors with on-off ratios of more than $10^6$ are readily available.

This review is organized in three sections as follows. In section 1, the state-of-the-art of the SVT is presented. It starts with a description of the device structure, the principle of operation, as well as essential fabrication technology. It then continues to cover two main topics in more detail: (i) the fundamental physics of the spin-dependent electronic transport and (ii) important device characteristics and performance. In section 2, the MTT is discussed. Since the principle of operation and the underlying physics is essentially the same as for the SVT, the focus is on those aspects of the MTT that are different from the SVT. In the third section, the use of hot-electron spin-filtering for spin-injection into a semiconductor is briefly discussed. Finally, a summary is given and some promising avenues for further research are identified.

1. Spin-valve transistor

1.1. Basic device structure and properties

The SVT was introduced in 1995 and is the first working hybrid device in which ferromagnets and semiconductors have been closely integrated, and both materials are essential in controlling the electrical transport through the device. The three-terminal device has the typical emitter/base/collector structure of a (bipolar) transistor, but is different in that the base region is metallic and contains at least two magnetic layers separated by a normal metal spacer (see figure 2). The two magnetic layers act as polarizer and analyser of electron spins, such that the relative orientation of the magnetization of the two layers determines the transmission of the base. The resulting salient feature of the SVT is that the collector current depends on the magnetic state of the base. This was first demonstrated by Monsma et al [21] in a SVT showing a huge change of 390% of the collector current in an applied magnetic field at low temperature. In 1998, the first device operating at room temperature was achieved, having a 15% effect in fields of a

![Figure 2. Basic layout of the SVT, showing the three terminal arrangement with semiconductor emitter (top), semiconductor collector (bottom), and the metallic base comprising two ferromagnetic thin layers separated by normal metals (middle).](image-url)
few kOe [24]. More recently, we succeeded in the reproducible fabrication of SVTs that exhibit magnetocurrent effects up to 400% at room temperature, and in small magnetic fields of only a few Oe [22, 23].

Unlike other spintronic devices, the SVT is based on the spin-dependent transport of non-equilibrium, so-called hot electrons, rather than Fermi electrons. In order to illustrate this and explain the principle of operation, let us consider the specific structure that was used in [22]. That particular SVT uses silicon as the semiconductor for the emitter and collector, and has a metallic base that contains a Ni80Fe20/Au/Co spin valve (see figure 3). At the interfaces between the metal base and the semiconductors, energy barriers (Schottky barriers) are formed [25]. These energy barriers prevent electrons with the Fermi energy from travelling through the structure. To obtain the desired high quality Schottky barrier with good rectifying behaviour and thermionic emission dominating, low doped Si (1–10 Ω cm) is used, and thin layers of, e.g. Pt and Au are incorporated at the emitter and collector side, respectively. These also serve to separate the magnetic layers from direct contact with Si.

The operation of the SVT is as follows. A current is established between the emitter and the base (the emitter current $I_E$), such that electrons are injected into the base, perpendicular to the layers of the spin valve. Since the injected electrons have to go over the Si/Pt Schottky barrier, they enter the base as non-equilibrium, hot electrons. The hot-electron energy is determined by the emitter Schottky barrier height, which is typically between 0.5 and 1 eV, depending on the metal/semiconductor combination [25]. As the hot electrons traverse the base, they are subjected to inelastic and elastic scattering, which changes their energy as well as their momentum distribution. Electrons are only able to enter the collector if they have retained sufficient energy to overcome the energy barrier at the collector side, which is chosen to be somewhat lower than the emitter barrier. Equally important, a hot electron can only enter the collector if its momentum matches with that of one of the available states in the collector semiconductor. The fraction of electrons that is collected, and thus the collector current $I_C$, depends sensitively on the scattering in the base, which is spin dependent when the base contains magnetic materials. The total scattering rate is controlled with an external applied magnetic field, which changes the relative magnetic alignment of the ferromagnetic Ni80Fe20 and Co layers in the base. This is illustrated in figure 4, where $I_C$ at room temperature is plotted as a function of magnetic field, for a transistor with a Si(100) emitter, a Si(111) collector and the following base: Pt(20 Å)/NiFe(30 Å)/Au(35 Å)/Co(30 Å)/Au(20 + 20 Å). At large applied fields, the two magnetic layers have their magnetization directions aligned parallel. This gives the largest collector current ($I_C = 11.2$ nA). When the magnetic field is reversed, the difference in switching fields of Co (22 Oe) and NiFe (5 Oe) creates a field region where the NiFe and Co magnetizations are antiparallel. In this state, the collector current is drastically reduced ($I_C^{AP} = 3.3$ nA). The magnetic response of the SVT, called the magnetocurrent (MC), is defined as the change in collector current normalized to the minimum value, i.e.

$$MC = \frac{I_C - I_C^{AP}}{I_C^{AP}},$$

where P and AP refer to the parallel and antiparallel magnetic state of the base spin valve, respectively. Thus, the relative magnetic response is MC = 240%, which is huge indeed.

It was shown [22, 26], that the collector current and the MC are virtually independent of a reverse bias voltage applied across the collector Schottky barrier, provided that the intrinsic ‘leakage’ current this induces in the collector diode is negligible compared to the hot-electron current. This is indeed the case, as can be seen in figure 5. A voltage between base and collector does not affect the hot-electron current because it does not significantly change the maximum of the Schottky barrier when measured with respect to the Fermi energy in the metal [25]. In other words, the energy barrier seen by hot electrons coming from the base is hardly changed. Similarly, a change of the emitter to base voltage, or equivalently the emitter current, does not affect the energy at which hot electrons are injected into the base. Also here, the applied voltage hardly modifies the maximum of the emitter potential barrier. The result is that the collector current is simply linearly proportional to the emitter current, as shown in figure 6.

It should be noted here that huge or even colossal magnetic response has been observed in other materials and devices. The uniqueness of the SVT is that the huge relative magnetic effect is obtained at room temperature, and that only small magnetic fields of a few Oe are required. The combination of these

![Figure 3. Schematic layout and energy band diagram of a SVT, showing the semiconductor emitter (left) and collector (right), and the metallic base comprising a spin valve (middle). Also depicted is the stream of electrons that is injected into the spin valve base above the Fermi energy.](image)

![Figure 4. Collector current versus applied magnetic field for a SVT with Si(100) emitter, Si(111) collector, and the following base: Pt(20 Å)/NiFe(30 Å)/Au(35 Å)/Co(30 Å)/Au(20 + 20 Å). $I_E = 2$ mA, $V_{BC} = 0$ and $T = 295$ K.](image)
three features is what makes it attractive. Furthermore, the results are reproducible and the properties of the device can be manipulated to a certain extent by controlling the thickness of layers, the type of materials, etc. Nevertheless, a particular point that needs attention is the absolute value of the output current. In the above example, an emitter current of 2 mA was used and the corresponding transfer ratio, defined as \( \alpha = I_C/I_E \), is thus of the order of \( 10^{-6} \). In this context, it is important to remember that the term ‘transistor’ should not be confused with ‘amplifier’. The word transistor was chosen in analogy with the metal base transistor (MBT) [27–29], a device that is similar to the SVT except that the base contains only non-magnetic metals. These structures were studied extensively in the 1960s and 1970s, from which it became clear that a metallic base has too little transmission to support amplification [27–29]. Hence, the SVT should not be judged on its (lack of) amplification. Rather, it should be viewed as a device with a magnetic field dependent electrical output, which is the basic functionality one needs for a magnetic field sensor or a magnetic memory element. Although current gain is thus not required for most applications, a small absolute current is a disadvantage. This issue and the progress made so far is addressed in one of the subsequent sections.

Before discussing the device in more detail, the essential role of the semiconductors in providing potential energy barriers must be emphasized. The energy barrier at the emitter side is needed to create injection of hot electrons, such that transport is governed by non-equilibrium processes. The collector energy barrier acts as an energy and momentum filter, allowing only a fraction of the hot electrons to pass into the collector, and reflecting the rest. The collector Schottky barrier (Au/Si in the above example) selects only on the basis of energy and momentum, but not on the spin of the incoming hot electrons. In some sense, the role of spin is thus indirect as it is merely used to manipulate the energy and momentum distribution of the hot electrons during their motion through the base. This is essential though, as spin couples to an external applied magnetic field and is our handle to the outside world. As will be discussed in detail below, the origin of the large magnetic sensitivity of the SVT is the non-equilibrium nature of the transport that gives an exponential decay of base transmission, together with the strong spin dependence of the elastic and inelastic scattering parameters in ferromagnetic materials.

1.2. Fabrication: vacuum metal bonding

For the correct operation of the SVT, it is crucial to have high-quality Schottky contacts that form defect-free energy barriers. This requires device quality semiconductor material. Since there is no known method of growing crystalline Si with low defect density on top of a thin metal film, an alternative fabrication technology based on metal bonding was developed [30]. The basic idea, illustrated in figure 7, is as amazing as it is essential for the SVT. In ultrahigh vacuum, we first deposit part of the base metal layers onto the emitter Si wafer, using a shutter to prevent deposition on the collector Si wafer. Then the shutter is opened, metal is deposited onto both wafers simultaneously, during which the surfaces of the metal coated wafers are brought into contact. A metal–metal bond is formed that bonds the two wafers together. The whole process of metal deposition and bonding is done \textit{in situ} under ultrahigh vacuum conditions in a molecular beam epitaxy system. This yields clean and incredibly strong metallic bonding.

As described previously [30, 31], a kind of recrystallization occurs when the two metal surfaces are brought into contact. This can repair structural defects at the bonding interface, provided the two surfaces are not too rough and atomic diffusion is sufficiently large. The driving force for the formation
of the chemical bonds is the gain in free energy when two surfaces are combined to form a bulk crystal structure. The quality, strength and reliability of the metal bond therefore depends on the type of metal used, as well as on the particular film topography (roughness, etc).

To study the bonding interface, we performed metal bonding in structures that are simpler than the SVT in that they contain only one type of metal, instead of the complete spin valve. The bonding interfaces were then investigated by transmission electron microscopy (TEM). Two examples are shown in figure 8, where two Si substrates are bonded together by two 5 nm thick Co layers (left panel) or two Au layers (right panel). For the Co structure, one can clearly recognize a lighter band that indicates a high density of structural defects at the bonded interface (denoted by the dashed arrow). In contrast, for the case of Au one recognizes parallel atomic (111) planes running across the complete Au layer and no clear bonding interface is visible. Although we have mostly used Au, good results have also been obtained with other metals such as Cu and Ti [30, 31]. For practical reasons related to the mechanical design of the ultrahigh vacuum bond robot, the bonding is typically done at room temperature with 1 cm² pieces of Si. Special care has to be taken to work in dust-free conditions and remove edge particulates after sawing of the wafers. We note that bonding is not limited to Si, but is a versatile technique that allows one to join many kinds of semiconductors. For example, SVTs with GaAs as the emitter and Si as collector have been successfully fabricated using vacuum metal bonding [32].

After metal bonding, the structures are further processed into smaller devices using standard photolithography and a series of dry and wet etching steps. The first step is to pattern the emitters into squares ranging from $350 \times 350 \mu m^2$ to $1 \times 1 mm^2$, using lithography and wet etching. The most convenient and reproducible way to do this uses a silicon-on-insulator (SOI) wafer for the emitter. After bonding, the handle wafer of the SOI wafer is first etched away, using the buried oxide as the etch stop. Subsequently, the oxide is etched away, leaving a thin Si layer of homogeneous and well-defined thickness. This layer, $3 \mu m$ thick, has a doping profile already built in. The front side is low doped (concentration $10^{15} \text{cm}^{-3}$, n-type), to form a proper Schottky barrier with the metal base. The back side is highly doped such that a good Ohmic contact is obtained by deposition of Cr and Au. The next step is patterning of the metal base into slightly larger rectangles using ion milling. A drawback is that this creates damage in the Si collector, which results in large leakage currents along the edges of the metal/collector Schottky contacts. This is detrimental as it reduces the MC by adding only to the total collector current (denominator in equation (1)). The edge leakage currents are found to depend strongly on temperature, and, as shown previously [22, 26], may prevent a large magnetic response being obtained at room temperature. We therefore introduced an additional process step to remove the damaged collector Si with a last wet chemical etch. Collector Schottky diodes with reverse bias currents of the order of 0.1 nA can thus be obtained.

The final standard chip consists of a total of 52 SVTs with different sizes (see figure 9). Electrical connections to these relatively large SVTs are made by ultrasonic wire bonding. This is however no longer feasible if smaller transistors are to be made. A modified fabrication process was therefore developed for the miniaturization of the device [33, 34], for which the use of SOI emitter wafers is essential. Devices with interconnect metallization and bondpads have been made so far with sizes down to $10 \times 10 \mu m^2$, and with properties comparable to the larger SVTs for which results are presented here.

1.3. Physical basis: hot-electron spin-transport

Unlike other spintronic devices, the SVT is based on the spin-dependent transport of hot electrons, rather than Fermi electrons. Transport at the Fermi energy has been widely studied in connection with (giant-) magnetoresistance effects and it is well-established that conduction in ferromagnets and their multilayers is dependent on the spin of the electrons [35]. However, hot-electron transport is distinctly different from ordinary transport at the Fermi energy. Therefore, to understand the operation of the SVT and avoid erroneous interpretation of results, it is best to 'forget' about what we know about GMR and spin-dependent scattering in spin valves.

It has been known for quite some time that scattering rates of hot electrons in ferromagnetic materials are spin dependent.
Early work on hot-electron scattering at energies of about 5 eV employed spin-polarized photoemission from overlayer structures [36, 37]. This and also later experiments [38, 39] at energies as low as 1.5 eV showed unambiguously that the inelastic mean free path of hot electrons is spin dependent in ferromagnets. More precisely, experiments carried out thus far have always found that the inelastic mean free path of hot electrons is shorter for the minority spin electrons [40]. A common interpretation is that this originates from the difference in the number of unoccupied states for the hot electron to scatter into, assuming electron–hole pair excitations to be the dominant scattering mechanism. Scattering processes have also been investigated using transmission through freestanding magnetic thin film foils [41–44], as well as by a time-resolved, two-photon photoemission (2PPE) experiment [45] that directly probed the inelastic lifetime for majority and minority spin.

Spin-dependent scattering of hot electrons is essential in a variety of spin-polarized electron spectroscopies that are widely used to examine magnetic materials. It leads to spin filtering and in magnetic multilayers it gives rise to phenomena such as the hot-electron spin-valve effect, first observed for secondary electrons [46]. This is employed in the magnetic version of ballistic electron emission microscopy (BEEM) [47], which enables magnetic imaging with nanometer resolution [48, 49]. With the introduction of the SVT [21, 24], which enables magnetic imaging with nanometer resolution [48, 49], a spin-polarized hot-electron transport was implemented in a solid-state electronic device. A thorough understanding of the spin-dependent scattering mechanisms of hot electrons is paramount to the further development of this type of device. Interestingly, the SVT itself has opened up a new route to study spin-dependent scattering processes of hot electrons, extending experiments to lower energy in the range between 0.5 and 1.5 eV. Moreover, for a solid-state device like the SVT it is rather easy to vary experimental parameters like temperature or apply large magnetic fields. This has led to some new insights into the origin of the spin-dependent scattering processes, as discussed below.

It is now well established that the difference between $T_p^C$ and $T_{AP}^C$ originates from the spin asymmetry in hot-electron attenuation length, leading to the dominant transmission of majority spins in each magnetic layer. This is schematically illustrated in figure 10. We write:

$$I_C^P \propto T_{NiFe}^M T_{Co}^M + T_{NiFe}^m T_{Co}^m$$  \hspace{1cm} (2)

$$I_C^{AP} \propto T_{NiFe}^M T_{Co}^M + T_{NiFe}^m T_{Co}^m$$  \hspace{1cm} (3)

where $T^M$ and $T^m$ are the transmission of majority (M) and minority (m) spin hot electrons, given by

$$T_{fi}^M = \Gamma_{in}^M \exp \left( -\frac{\lambda_{fi}^M}{\lambda_{in}^M} \right) \Gamma_{out}^M,$$  \hspace{1cm} (4)

$$T_{fi}^m = \Gamma_{in}^m \exp \left( -\frac{\lambda_{fi}^m}{\lambda_{in}^m} \right) \Gamma_{out}^m.$$  \hspace{1cm} (5)

Here, $\lambda_{fi}$ is the film thickness, $\Gamma_{in}$ and $\Gamma_{out}$ the hot-electron attenuation lengths for majority and minority spin, and $\Gamma_{fi}$ denotes the ferromagnet. The factors $\Gamma_{out}$ denote the transmission across the interface of the ferromagnet with the normal metal at each side, which can be spin dependent due to the mismatch in the bandstructure of the magnetic and non-magnetic metals. Thus, the interfacial attenuation is generally dependent on the combination of ferromagnet and normal metal. Now, an initially unpolarized current acquires, after transmission of a film $fm$, a spin polarization $P$ given by

$$P_{fm} = \frac{T_{fm}^M - T_{fm}^m}{T_{fm}^M + T_{fm}^m}.$$  \hspace{1cm} (6)

This is related directly to the magnetocurrent $MC$ via:

$$MC = \frac{I_C^P - I_C^{AP}}{I_C^P} = \frac{2 P_{NiFe} P_{Co}}{1 - P_{NiFe} P_{Co}}.$$  \hspace{1cm} (7)

In general, the attenuation is caused by a combination of inelastic and elastic scattering processes. It is important to realize that while the scattering parameters are uniquely defined for a given ferromagnet and hot-electron energy, the attenuation length is not [50]. The attenuation length depends not only on the scattering parameters, but to some extent also on the device geometry and the energy and momentum selection applied by the collector Schottky barrier. For example, in the hypothetical case that a collector is used which passes hot electrons regardless of their momentum, elastic scattering in the base metals would not directly cause attenuation, and the attenuation length would be governed mainly by inelastic processes.

Let us address the origin of the large magnetocurrent effect and consider the factors that distinguish the SVT from magnetoresistive structures such as a magnetic multilayer. First of all, transport is by hot electrons, which probe a different portion of the band structure than Fermi electrons, while the dominant scattering mechanisms are different and include inelastic scattering such as by electron–hole pair excitations. Second, the resistance of metals is governed by the mean free path, and the resistivity is inversely proportional to the mean free path. In contrast, the hot-electron transport in the SVT is a

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1 Some indirect effect of elastic scattering still occurs because it changes the path length in the base and thereby increase the probability for an inelastic scattering event.
non-equilibrium phenomenon where the current transmission depends exponentially on the base thickness, and the parameter that controls the decay is the attenuation length. As stated above, this is related to the mean free path, but also depends on the energy and momentum selection at the collector barrier. Furthermore, the relative magnetic effect (the MC) of the SVT is rather insensitive to scattering that carries no spin dependence, as it attenuates the current for both spins with an equal factor, leaving the ratio unchanged. A final point is that transport is largely perpendicular to the magnetic layers, such that the electrons have to cross all the interfaces. All these factors together facilitate the large effects.

### 1.4. Interface versus volume scattering

The reproducible fabrication of the SVT permitted a detailed study of the spin-dependent transmission of low energy hot electrons through ferromagnetic films and trilayer structures. A first question that was addressed is, what are the typical attenuation lengths \( \lambda_m \) and \( \lambda_M \) for majority and minority spin hot electrons in a ferromagnetic metal? Another interesting question is whether the dominant scattering occurs within the volume of the magnetic layers or at the interfaces with the non-magnetic layers. With respect to spin-dependent scattering, one may expect volume scattering to dominate as its contribution grows exponentially with ferromagnetic film thickness and will quickly outweigh the (constant) interfacial scattering, one may expect volume scattering to dominate as its contribution grows exponentially with ferromagnetic film thickness and will quickly outweigh the (constant) interfacial contribution (see equations (4) and (5)). This is indeed what experiments with the SVT suggest [51,52]. Such experiments involve fabrication of a series of SVTs that are identical except for the thickness of one of the ferromagnetic layers of the base.

Let us discuss the particular example of a transistor with the following base structure Pt(30 Å)/Ni_{80}Fe_{20}(x)/Au(44 Å)/Co(30 Å)/Au(44 Å), where the thickness \( x \) of the Ni_{80}Fe_{20} layer was varied. Figure 11 shows \( I_P \) and \( I_{AP} \) (top panel) and MC (bottom panel) at \( T = 100 \text{ K} \), as a function of the thickness \( x \) of the Ni_{80}Fe_{20} layer, keeping the Co layer thickness constant at 30 Å. For the parallel magnetic state of the spin valve, minority spins are strongly attenuated in both magnetic layers, and only majority spins contribute to \( I_P \). From the exponential decay of \( I_P \) with thickness, we thus deduce a volume attenuation length of 43 ± 3 Å for majority hot spins in Ni_{80}Fe_{20}.

The interface attenuation for majority spins is extracted by comparing the \( I_C \) value of 169 nA for the device with \( x = 0 \), to the value of 78 nA obtained by extrapolation of the \( I_P \) data to zero thickness. The latter value is a factor of 2.2 lower, which is due to attenuation by the extra interface that is created when a Ni_{80}Fe_{20} layer is inserted between the Pt and the Au layer of the base (more precisely, a factor of 2.2 represents the difference between the original Pt/Au interface, and the two new Pt/Ni_{80}Fe_{20} and Ni_{80}Fe_{20}/Au interfaces). Note that this is the attenuation for the majority spins, as the minority spins are filtered out by the 30 Å Co layer, irrespective of their scattering at the Ni_{80}Fe_{20} interfaces. The interfacial attenuation is a combination of the mismatch of the electronic states at both sides of the interface, and elastic scattering due to interface disorder, defects, etc.

Information on the spin dependence of the volume and interface attenuation is contained in the difference between \( I_P \) and \( I_{AP} \). Let us first compare the data with what is expected for two extreme cases, using equations (2)–(5) for the collector current. In the first case, represented by the dashed lines in figure 11, all the spin dependence is assumed to arise from the interfacial scattering, while identical volume attenuation lengths are used for both spins. In that case the MC would be independent of Ni_{80}Fe_{20} thickness, and \( I_P \) and \( I_{AP} \) decay with \( x \) with the same slope. The experimental data clearly deviates from this behaviour at small thickness. In contrast, the data agrees very well with the other extreme case (solid lines), in which the interface scattering is not spin dependent, and the only spin-asymmetry is that of the volume attenuation lengths. In this situation we expect that \( I_P \) and \( I_{AP} \) approach each other as \( x \) goes to zero, such that the MC tends to zero. Note that for large \( x \), filtering of minority spins is complete and the decay of both \( I_P \) and \( I_{AP} \) is determined by the majority spin attenuation length.

More precise analysis, taking into account the experimental error margins, shows that some weak spin dependence of the interface scattering cannot be excluded [51]. For the ratio of the interface attenuation factor \( \Gamma \) for minority and majority spin we obtain \( \Gamma^{\uparrow}/\Gamma^{\downarrow} = 0.8 \pm 0.2 \), while volume attenuation lengths are \( \lambda_{NiFe} = 10 \pm 2 \text{ Å for minority spin} \) and \( \lambda_{NiFe} = 43 \pm 3 \text{ Å for the majority spin (at 0.9 eV)} \). With these parameters, a 30 Å thick Ni_{80}Fe_{20} film attenuates minority spins a factor of 10 more strongly than the majority spins. Volume scattering is thus by far the dominant contribution to the spin-dependence of the hot-electron attenuation, mainly because it depends exponentially on the layer thickness.
1.5. Elastic versus inelastic

In hot-electron scattering, it is obvious that inelastic scattering plays a role, as it leads to energy loss and thus makes the electrons ‘less hot’. However, elastic scattering cannot be neglected a priori. The point is illustrated best by asking what we would expect if there were no elastic scattering whatsoever, neither in the volume of the layers, nor at interfaces. Can we explain the basic electrical characteristics of the SVT using only inelastic scattering? A simple estimate based on exponential decay of the transmission of a particular layer is illuminating. It is well-known from MBTs [27–29] and BEEM [53, 54] that scattering lengths for non-magnetic materials such as Au or Cu are quite long, at least 100 Å. In the SVT, the thickness of these normal metal layers is typically only 30–60 Å such that inelastic attenuation in these layers can be neglected. In order to satisfy the experimental observations (a typical MC of 400% at low T and an overall transmission of 10⁻¹ for SVTs with ferromagnetic layers of 30 Å (see section 1.7)), one would require inelastic attenuation lengths of 7 Å and 4.5 Å for majority and minority spins, respectively.

Such short inelastic scattering lengths appear unrealistic. Moreover, they are at odds with a recent direct measurement [45] of the inelastic lifetime of hot electrons by 2PPE, which yielded a lifetime of 8 fs and 4 fs for majority and minority spin hot electrons, respectively, in Co at an energy of about 1 eV. With a typical electron velocity of 2 × 10⁶ m s⁻¹, this translates into inelastic scattering lengths of about 160 and 80 Å for majority and minority spins in Co. These values are much larger than those quoted above. For minority spins, a significant density of d-states is still present at energies of 1 eV above the Fermi energy, and one could argue that the electron velocity is much smaller than assumed above. Yet, such an argument cannot be made for majority spins, for which only highly mobile states with s or p character are available. It is possible that there is a principle reason why inelastic lifetimes determined by the optical technique of 2PPE are much larger than those observed in electronic transport. This may be [55] due to diffusion of electrons away from the optical spot, due to matrix elements for optical transitions, or due to a different energy resolution of the techniques.

Attenuation lengths for several ferromagnetic materials in the SVT [51], the MTT [56], and in magnetic BEEM [49] all consistently give minority spin attenuation lengths around 10 Å, and majority scattering lengths in the range of 25–60 Å. These measured values are significantly larger than values of 4.5 and 7 Å estimated above for the case in which only inelastic scattering is present, but smaller than what is derived from the 2PPE results. A consistent picture can only be made if elastic scattering is included. First, transport attenuation lengths that are determined by elastic and inelastic processes are naturally shorter than predicted from 2PPE, which probes only the inelastic scattering lifetime. Second, when measured attenuation lengths are used to calculate the total volume attenuation of the base layers, one arrives at a transmission of about 0.01. Since observed base transfer ratios are at least 2 orders of magnitude smaller, a significant fraction of the attenuation must be attributed to scattering at interfaces, which is primarily elastic.

Finally, we mention that attributing the measured transport attenuation length solely to inelastic scattering is inconsistent with the general trend observed in SVT [57], MTT [56] and BEEM data [49] that the overall transmission increases with increasing hot-electron energy (see also section 2 on the MTT). For inelastic scattering, the opposite should happen, since according to Fermi’s golden rule the lifetime is reduced when at higher energy more final states become available to scatter into. This is a most compelling argument and we conclude that the experimental data cannot be explained if only inelastic processes are considered. The picture of purely ballistic transmission, thus, seems rather inappropriate.

Experimental evidence for the importance of elastic scattering comes from the explicit observation of attenuation due to interfaces (figure 11 in the previous section) and from the direct observation of spin-orbit scattering [58]. In addition, we have found that the magnitude of the collector current is insensitive to crystal orientation (100) or (111) of the Si collector (see data in table 1). Yet, a most instructive way to examine the role of elastic scattering is to make use of phenomenological model calculations, in which scattering parameters can be systematically varied and the effect on magnetotransport can be studied. The model and the calculation procedure are described in detail in [50]. Here, we merely discuss an example that illustrates the role of interfacial elastic scattering.

For the case in which elastic interface scattering was set to zero, the top panel of figure 12 shows how the calculated collectable current decays as the hot electrons traverse the metal base. Here, ‘collectable’ refers to those electrons that still satisfy the energy and momentum constraints required for transmission into the Si collector [50]. We observe an exponential attenuation of the current in the volume of each layer, with the strongest decay in each of the ferromagnetic layers. The bottom panel of figure 12 shows the results if the interface diffusivity is set to 0.9 for each interface in the base (including the metal semiconductor interfaces at the emitter and collector side). This corresponds to 90% of the electrons incident on an interface being scattered elastically, with scattered electrons distributed isotropically over all angular directions. One notes a pronounced drop in the calculated collectable current at each interface. This adds up to a large reduction of the overall base transmission. For the emitter current of 2 mA used, the total transmitted current Iₑ is only 8 × 10⁻⁹ A for the strong interface scattering case, as opposed to 8 × 10⁻⁶ A without interface scattering. The difference is as much as three orders of magnitude. Thus, elastic scattering at interfaces, for example, due to disorder, can cause a significant reduction of the collector current. Elastic scattering therefore plays a perhaps surprisingly important role in hot-electron transport.

Table 1. Transfer ratio α for several Si/Pt/Au/Si transistors with different crystal orientation of the collector. The Schottky barrier height Φ and ideality factor n are also given.

<table>
<thead>
<tr>
<th>Emitter</th>
<th>Base (Å)</th>
<th>Collector</th>
<th>Φₑ (eV)</th>
<th>n</th>
<th>α₁⁴³⁻⁻¹⁰⁻⁴</th>
</tr>
</thead>
<tbody>
<tr>
<td>Si(100)</td>
<td>0.91 1.01</td>
<td>40Pt/40Au</td>
<td>Si(111)</td>
<td>0.80</td>
<td>1.20</td>
</tr>
<tr>
<td>SOI(100)</td>
<td>0.85 1.07</td>
<td>40Pt/40Au</td>
<td>Si(111)</td>
<td>0.76</td>
<td>1.20</td>
</tr>
<tr>
<td>SOI(100)</td>
<td>0.88 1.02</td>
<td>40Pt/40Au</td>
<td>Si(100)</td>
<td>0.77</td>
<td>1.15</td>
</tr>
</tbody>
</table>
1.6. Temperature and spin waves

For ferromagnets, experiments carried out this far have always found that the inelastic mean free path of hot electrons is shorter for the minority spin electrons [40]. A common interpretation is that this originates from the difference in the number of unoccupied states for the hot electron to scatter into, assuming excitation of electron–hole pairs to be the dominant scattering mechanism. However, recent experimental [59] and theoretical work [60, 61] suggests that scattering by spin-wave excitations may also contribute to the spin dependence of the mean free path. The spin asymmetry is created by spontaneous spin-wave emission, which, due to conservation of angular momentum, is allowed only for minority spins. Spontaneous spin-wave emission is not dependent on temperature ($T$) and occurs even at $T = 0$. It is calculated [60, 61] to dominate at electron energies below $\approx 3$ eV, and is found to be the major source of small energy losses [59]. Besides spontaneous emission, there is also a thermal component which involves spin-wave absorption as well as emission. Information on these contributions is obtained most directly by varying $T$. While spectroscopic measurements have been performed, the $T$-dependence had so far not been explored. Using the SVT, we have obtained the first measurements of spin-dependent hot-electron scattering as a function of temperature [23]. We found that thermal spin waves have a noticeable effect on spin-transport of hot electrons, in particular by introducing spin-mixing.

Figure 13 shows the typical variation of the collector current with applied magnetic field at two temperatures (80 and 290 K). At high fields the magnetizations of the Co and Ni$_{80}$Fe$_{20}$ films in the spin-valve base are parallel. When one of the magnetizations is reversed, a clear antiparallel (AP) magnetization alignment is obtained. At $T = 80$ K, the magnetocurrent $MC = 560\%$, while a huge effect of 350% still remains at room temperature. Similar results were reproducibly obtained.

In figure 14, the temperature variation of $I_C^P$ and $I_C^{AP}$ (top panel) and the magnetocurrent (bottom panel) for the same SVT as in figure 13 are displayed. For the parallel case, the collector current first increases with $T$, but decreases above 200 K. In contrast, $I_C^{AP}$ increases over the whole temperature range. The resulting decay of the MC is relatively weak, especially below 200 K. The conventional current-in-plane magnetoresistance (CIP-MR) measured on an identical spin valve is also included. The CIP-MR is not only orders of magnitude smaller, but also exhibits a stronger thermal decrease that is linear, as also found by others [62].
To understand the data in figure 14 we consider that thermal scattering may change the MC in two ways. First, scattering may attenuate, i.e. remove electrons from the collected current in a spin-dependent fashion. The second possibility is spin-mixing, in which electrons are scattered into the other spin channel by a spin-flip process, without being removed from the collected current. We have shown [23] that when we neglect possible thermally-induced spin-mixing, the MC does not decay with $T$. Thus, the decay of MC is attributed to thermally-induced spin-flip scattering, causing mixing of the two spin channels. This is consistent with the observation (figure 2) that the current for the parallel state goes down above 200 K, while the current for the antiparallel case continues to go up. Since exchange scattering by paramagnetic impurities and spin-orbit scattering have negligible $T$-dependence in our experimental range [63], we consider spin-flip scattering by thermal spin waves. We stress that spin-mixing only results for those scattering events in which the energy and/or momentum transferred to the spin-wave is such that after scattering, the hot electron direction of motion is small, the deflection is weak. We write $\lambda_{e-h}^{2}$ for the mean free path associated with electron–hole pair excitations, $\lambda_{el}$ and $\lambda_{ph}$ of the attenuation lengths associated with elastic scattering and phonon scattering, respectively. Also included are attenuation lengths $\lambda_{TSW \ absorb}^{m}$ and $\lambda_{TSW \ emit}^{m}$ for absorption and emission of thermal spin waves, respectively, as well as a term $\lambda_{SSW \ emit}^{m}$ due to spontaneous emission of spin waves. Due to the conservation of angular momentum, only majority-spin hot electrons can absorb spin waves, whereas (spontaneous and thermal) emission is allowed only for minority spins. Thus, the overall rate of spin wave scattering has a spin asymmetry due to $\lambda_{SSW \ emit}^{m}$.

In order to extract the thermal component of the hot-electron attenuation, we have measured $I_{C}^{p}$ as a function of temperature [51]. We divide out all temperature independent attenuation factors by plotting the normalized collector current $I_{C}^{N} = I_{C}^{p}(T)/I_{C}^{p}(100 \text{ K})$ (see figure 15). The top panel shows data for transistors with different Ni$_{80}$Fe$_{20}$ thickness and constant Co thickness. Starting at 100 K the collector current first goes up slightly as seen before, and then is reduced significantly towards room temperature. Moreover, the variation with $T$ is more pronounced at higher Ni$_{80}$Fe$_{20}$ thickness, implying that the additional attenuation at higher $T$ is due to a thermal volume scattering process. The same behaviour is observed when the thickness of the Co layer is varied (bottom panel of figure 15). However, the $T$-variation for Co has weaker dependence on layer thickness than for Ni$_{80}$Fe$_{20}$, which shows that the thermal attenuation is stronger in Ni$_{80}$Fe$_{20}$ than it is in Co. This is consistent with attenuation due to scattering on thermal spin-waves [23], as the Curie temperature of Co (1388 K) is larger than that of Ni$_{80}$Fe$_{20}$ (873 K).

Using the procedure described in [51] one can extract the attenuation length $\lambda_{TSW}^{m}(T)$ due to thermal spin waves. The extracted thermal attenuation lengths for Ni$_{80}$Fe$_{20}$ and Co versus $T$ are shown in figure 16. For Ni$_{80}$Fe$_{20}$ two sets of data are shown, one for the SRT, the other labelled AMT is extracted from another set of transistors with

![Figure 15](image-url)
Pt(20 Å)/Ni80Fe20/Au(44 Å) base. As noted before, $\lambda_{\text{TSW}}(T)$ is shorter for Ni80Fe20 than for Co. Interestingly, $\lambda_{\text{TSW}}(T)$ is much shorter than hitherto assumed, with room temperature values of 130 ± 20 Å for Ni80Fe20 and 270 ± 40 Å for Co. Especially for Ni80Fe20 this is only three times larger than the majority spin attenuation length at low T. Hence, this shows convincingly that hot-electron attenuation lengths are dependent on T. For instance, the addition of thermal spin wave scattering with a length scale of 130 Å reduces the attenuation length from 43 Å at 100 K, to a significantly lower value of $(1/43 + 1/130) = 32$ Å at room temperature.

The inset of figure 16 shows the thermal variation of the interface attenuation $\Gamma(T)$ for Ni80Fe20. Data is obtained by dividing the $P^2$ curve for certain Ni80Fe20 thickness by the curve for $x = 0$, and using the attenuation lengths of figure 16 to remove the volume scattering part. We find only a slight change of about 5% in the interface attenuation between 100 and 300 K. This shows that thermal spin wave scattering is primarily a volume scattering process.

The above results suggest the importance of spin waves for the spin asymmetry of hot-electron transmission. As shown in [60, 61], the asymmetry is created by the T-independent contribution of spontaneous spin wave emission, which is only allowed for minority spins. It has so far not been possible to isolate spontaneous spin wave emission and measure the corresponding scattering length. However, the above quantification of the thermal spin wave contribution allows us to estimate the attenuation due to spontaneous emission, which is expected to be significantly stronger than the thermal scattering rate. This is because thermal spin waves occupy only a small fraction of the spin-wave phase space up to energies of the order of kT, while for spontaneous spin wave emission, the complete phase space up to the hot-electron energy ($\approx 0.9$ eV $\gg$ kT) is available at all temperatures. The strength of the observed attenuation due to thermal spin waves is thus indirect evidence for the importance of spontaneous spin-wave emission. If we crudely estimate the spontaneous emission rate to be about an order of magnitude larger than the thermal emission rate, we obtain an attenuation length for spontaneous emission that is close to the measured minority-spin attenuation length (10 Å for Ni80Fe20). This strongly suggests that the minority spin attenuation is dominated by spontaneous spin-wave emission, as theory predicts [60,61]. Since the process cannot contribute to attenuation of majority hot spins, we conclude that the spin asymmetry of the attenuation length may well be due to spontaneous spin-wave emission, instead of the spin-dependent rate of electron–hole pair generation that arises from the exchange split bandstructure.

The notion of spontaneous spin-wave emission as the dominant source of spin-asymmetry explains the absence of bandstructure features in spectroscopic data. In particular, it explains the puzzling observation [55] by time-resolved 2PPE, that the inelastic lifetime for minority spins is shorter than the majority spin lifetime, not only in Co and Ni, but also in Fe. For Fe, a reversed spin-asymmetry at low energy ($<1$ eV) is expected from the bandstructure [55]. However, spontaneous spin wave emission would, as it is forbidden for majority spins, always result in a larger lifetime for majority spins, as observed in the 2PPE experiment.

1.7. Device output and transfer ratio

In close harmony with the more fundamental research described in the preceding sections, quite some effort has been devoted to improve one of the most important device parameters, namely, the output current level. Surely, the prospects of application of the SVT will be greatly enhanced if the collector currents can be raised to well above the nanoampere range, such that one can capitalize on the huge magnetic response of the device. Fortunately, significant progress has been made in recent years [57], and there is still plenty of scope for further improvement.

From the basic relation $I_C = \alpha I_E$ we see that we can increase $I_C$ by either increasing the emitter current ($I_E$), or the transfer ratio ($\alpha = I_C/I_E$) of the transistor base. However, we know that the emitter current has an upper limit imposed by the breakdown of the device [34]. Therefore, we should optimize $\alpha$. Parenthetically, the most fruitful way to achieve this is to change the non-magnetic layers of the base. In a typical SVT structure of Si(100)/NiFe/Ni82Fe18/Au/Co/NiFe/Si(100), thin normal metal layers NiFe and NiFe are incorporated at the emitter and collector side of the structure, respectively. The primary function of these layers is to create high quality Schottky barrier diodes with thermionic emission dominating, to ensure efficient hot-electron injection from the emitter and low reverse bias leakage of the collector diode. From a magnetic point of view these layers are not active, providing some flexibility in device design. Below we show that this can be exploited [57] to enhance the transfer ratio in two ways: (i) controlling the difference in height of the emitter and collector Schottky barrier, and (ii) by carefully selecting materials with large scattering lengths for hot electrons.

We start by investigating the role of the Schottky barrier height of emitter ($\Phi_E$) and collector ($\Phi_C$) on the transfer ratio. The Schottky barrier heights in the SVT are determined by the choice of semiconductor, silicon in our case, and the non-magnetic metal. We have systematically varied the Schottky
barrier metal in two series of SVTs [57]. In the first series, denoted as Pt series, NM_E is Pt and NM_C is either Pt, Au or Cu. In the second series, denoted as Au series, NM_E is Au and NM_C is either Au or Cu. The properties of the two series of transistors are summarized in table 2.

In figure 17, we show the parallel collector current versus Schottky barrier height difference (\(\Delta \Phi_B = \Phi_B - \Phi_C\)) of the two series. Both the Pt and Au series show an increase in collector current for larger \(\Delta \Phi_B\). The increase is not related to a change in hot-electron energy, since in each series the emitter Schottky barrier is fixed, and only the collector barrier is lowered as one goes from Pt to Au to Cu. The same studies have been performed on MBTs in which the spin valve has been omitted from the base. The structure is thus Si(100)/NM_E/NM_C/Si(100), where NM_E and NM_C perform the same function and are made of the same materials as used for the SVTs in table 1. The results (not shown) yield the same trend for the transfer ratio versus \(\Delta \Phi_B\) and material as for the SVT series.

The enhancement of \(I_C\) is due to the larger number of states available in the collector semiconductor. Electrons that are injected at an energy just above the collector barrier (small \(\Delta \Phi_B\)) arrive at the base/collector interface at energies just above the conduction band minimum of the collector Si. Near the bottom of the band the density of states is small and only few electrons can enter the base. The structure is thus Si(100)/NM_E/NM_C/Si(100), where NM_E and NM_C perform the same function and are made of the same materials as used for the SVTs in table 1. The results (not shown) yield the same trend for the transfer ratio versus \(\Delta \Phi_B\) and material as for the SVT series.

Table 2. Properties of SVTs with different combinations of emitter and collector barriers, where sv denotes a NiFe/Au/Co spin valve. \(T = 290\, \text{K}\).

<table>
<thead>
<tr>
<th>Base</th>
<th>(\Phi_B) (eV)</th>
<th>(\Phi_C) (eV)</th>
<th>(\Delta \Phi_B) (eV)</th>
<th>(\alpha) ((\times 10^{-4}))</th>
<th>MC (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Pt/sv/Pt</td>
<td>0.88</td>
<td>0.86</td>
<td>0.02</td>
<td>0.01</td>
<td>213</td>
</tr>
<tr>
<td>Pt/sv/Au</td>
<td>0.88</td>
<td>0.83</td>
<td>0.05</td>
<td>0.07</td>
<td>260</td>
</tr>
<tr>
<td>Pt/sv/Cu</td>
<td>0.87</td>
<td>0.61</td>
<td>0.26</td>
<td>1.06</td>
<td>218</td>
</tr>
<tr>
<td>Au/sv/Au</td>
<td>0.81</td>
<td>0.80</td>
<td>0.01</td>
<td>0.10</td>
<td>204</td>
</tr>
<tr>
<td>Au/sv/Cu</td>
<td>0.82</td>
<td>0.69</td>
<td>0.13</td>
<td>1.18</td>
<td>230</td>
</tr>
</tbody>
</table>

Figure 17. \(I_C^E\) for five types of SVTs versus \(\Delta \Phi_B\). The circles are the Pt series and the squares are the Au series. \(I_E = 2\, \text{mA}, \ T = 290\, \text{K}\) and sv denotes a NiFe/Au/Co spin valve.

Figure 18. Collector current versus applied magnetic field for a Si/Au(20 Å)/NiFe(30 Å)/Au(70 Å)/Co(30 Å)/Cu(40 Å)/Si transistor. \(I_E = 2\, \text{mA}\) and \(T = 290\, \text{K}\).
Figure 19. Parallel collector current versus Pt thickness $t_{Pt}$ for a Si/Pt/NiFe(30 Å)/Au(40 Å)/Co(30 Å)/Au(40 Å)/Si transistor. $I_E = 2$ mA and $T = 100$ K.

**Figure 19.** Parallel collector current versus Pt thickness $t_{Pt}$ for a Si/Pt/NiFe(30 Å)/Au(40 Å)/Co(30 Å)/Au(40 Å)/Si transistor.

Pt thickness, although the precise cause has not yet been identified.

The above example nicely illustrates the effect of the structural quality of the SVT on the transfer ratio, and thereby points to a promising route for further improvement. As shown in section 1.5, a significant part of the hot-electron attenuation is due to elastic scattering processes which likely originate from structural imperfections (defects, grain boundaries, stacking faults, disorder at interfaces, etc.). These are surely present in the polycrystalline, non-epitaxial structures we have studied so far. Going to epitaxial transistor structures is an approach we are currently exploring. Another option is to reduce the number of layers (and interfaces) in the base, and with it, the elastic scattering. This can be done in the slightly modified design of a MTT (see section 2 on the MTT). Finally, we note that some optimization of the thickness of the ferromagnetic layers is needed, since there is a trade-off between transfer ratio and MC that yields a maximum in absolute change of collector current ($\Delta I_C$) with magnetic field [24, 57].

In the preceding paragraphs we have focused on how to improve the transfer ratio, keeping the emitter current fixed at a somewhat arbitrary value of 2 mA. However, this is far below the breakdown current of the SVT [34]. Thus, the absolute value of the output collector current is by no means limited to the value of 0.2 $\mu$A shown in figure 18. To illustrate this, figure 20 shows how the collector current, transfer ratio and MC vary with emitter current, for a SVT with similar structure as in figure 18. We find that the output current of the SVT increases approximately linearly with emitter current, and large collector currents up to 50 $\mu$A are obtained. Interestingly, the increase of output current by several orders of magnitude is accompanied by a rather weak reduction of the MC, and even a slight increase of the transfer ratio. This convincingly demonstrates that the collector current is not intrinsically limited to small values in the nA regime, but can approach 0.1 mA. In the end, the output current that can be achieved depends on the maximum emitter current that can be tolerated for a given application. This is determined by extrinsic requirements such as power consumption, device dimensions since the breakdown emitter current depends on device size [34], and input impedance (for the maximum emitter current in figure 20, the applied emitter voltage was about 4 V, corresponding to an input impedance of 16 $\Omega$). This also implies that care has to be taken when comparing the properties of the SVT with similar devices such as the MTT, and correct conclusions can only be drawn if differences in the applied emitter current are properly considered [65].

### 1.8. Noise

An important parameter for any electronic device is the signal-to-noise ratio (SNR). Let us first examine the origin of the noise in the SVT. To characterize the noise behaviour over a wide range of collector current values, measurements have been performed on three different types of transistors [66]. Besides the SVT with Pt/NiFe/Au/Co/Au spin valve base, we also used a Pt/NiFe/Au MBT (with only a single magnetic layer), as well as the non-magnetic MBT with a Pt/Au base. The collector current noise spectrum showed only white noise in the frequency range of the measurement (10 Hz–1 kHz). The noise current spectral density is plotted as function of $I_C$ in the top panel of figure 21, where the $I_C$ for each of the three transistors was varied by changing $I_E$. The log–log plot yields a straight line with a slope equal to 1, which means a linear relationship between noise spectral density and collector current. This is expected when shot noise dominates. In fact, the experimental data are in excellent agreement with the calculated shot noise of $2qI_C$, as indicated by the solid line.
Calculated noise calculated from the measured magnetic field for the SVT. Solid lines in both panels are the shot transistors with three different base structures: Pt/NiFe/Au/Co/Au such a structure would drop significantly. The result would be a loss in SNR, because the collector current of the ferromagnetic layers (to 50 Å or more). In fact, the net straightforward to do so, one simply increases the thickness to create an SVT with MC values above 1000%, although it is intuitive, but is simply because the absolute value of the signal the full difference between $I_C$ and $P$ is already close to the maximum at an MC of several 100%. It is for this reason that we have never attempted to emphasize the need for further increasing the transfer ratio. One also notes that there is a significant gain in SNR between an MC of 10% and 100%, however, beyond 100% the gain levels off and only a marginal SNR increase is observed when the MC is further raised to 1000%. This may seem counter intuitive, but is simply because the absolute value of the signal $I_C^P - I_C^{NP}$ is already close to the maximum at an MC of several 100%. It is for this reason that we have never attempted to create an SVT with MC values above 1000%, although it is straightforward to do so, one simply increases the thickness of the ferromagnetic layers (to 50 Å or more). In fact, the net result would be a loss in SNR, because the collector current of such a structure would drop significantly.

For the SVT, the relation between noise and the magnetic state of the base was studied. As shown in figure 21, bottom panel, the noise varies with applied magnetic field very much like the collector current itself, and no extra noise appears in the field regions where the magnetization reversal occurs. The change of the noise with field is entirely due to the change of the collector current. The solid curve is the shot noise calculated from $2qI_C$, using the measured collector current versus field curve. Again, good agreement is obtained with the data. We conclude that shot noise dominates in the investigated current and frequency regime.

We can now evaluate the SNR of the device, using as the signal the full difference between $I_C^P$ and $I_C^{NP}$. The result is shown in figure 22, where SNR is plotted against $I_C$ for devices with different values of the MC percentage. It is evident that the SNR improves significantly at higher collector current, emphasizing the need for further increasing the transfer ratio. One also notes that there is a significant gain in SNR between an MC of 10% and 100%, however, beyond 100% the gain levels off and only a marginal SNR increase is observed when the MC is further raised to 1000%. This may seem counter intuitive, but is simply because the absolute value of the signal $I_C^P - I_C^{NP}$ is already close to the maximum at an MC of several 100%. It is for this reason that we have never attempted to create an SVT with MC values above 1000%, although it is straightforward to do so, one simply increases the thickness of the ferromagnetic layers (to 50 Å or more). In fact, the net result would be a loss in SNR, because the collector current of such a structure would drop significantly.

2. **Magnetic tunnel transistor**

The MTT is derived from the SVT and is so similar that it is often referred to as a tunnel SVT. The MTT differs in the structure of the emitter used to inject the hot electrons into the transistor base. Whereas the SVT uses a Schottky barrier, the MTT has a tunnel barrier as the emitter. In the first publication on the SVT in 1995 by Monsma et al [21], this alternative device geometry using a tunnel barrier structure as the emitter was already proposed. In 1997, Mizushima et al were the first to actually fabricate a MTT [67], and they own the patent [68].

The MTT device exists in two slightly different variations, as shown in figure 23. The design on the left has a thin insulating layer to separate the metal base from a non-magnetic metal emitter electrode. When a voltage $V_{EB}$ is applied across the tunnel barrier between emitter and base, unpolarized electrons are injected by tunnelling into the metal base, arriving at an energy $eV_{EB}$ above the Fermi level in the metal base. The rest is the same as in the SVT; spin-dependent transmission of the hot electrons through the two ferromagnetic thin films in the metal base, and collection across a Schottky barrier with energy and momentum selection.

The design on the right in figure 23 also has a tunnel barrier, but uses a ferromagnetic emitter electrode. Because the tunnel probability is spin-dependent in such a structure, the injected hot electrons are already spin polarized as they enter the transistor base. Therefore, the first polarizing magnetic layer in the base can now be omitted and only the second analyser magnetic layer needs to be retained. Again, the collector current is determined by spin-dependent transmission of the hot electrons through the ferromagnetic base and collection across a Schottky barrier with energy and momentum selection. Because only the emitter part is different, we can directly apply much of the knowledge obtained from studies of the SVT described in sections 1.3–1.8 of section 1. The exponential variation of the base transmission, the importance of elastic, inelastic and thermal scattering processes in the layers and at interfaces, and the conditions for transmission into the collector are essentially the same for MTT and SVT. As an example, hot-electron attenuation lengths measured with the MTT [56] are
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Figure 23. Energy diagrams of the MTT with a non-magnetic tunnel injector (left) or a ferromagnetic tunnel injector (right) to generate the hot-electron current.

comparable to that previously found for the SVT [51] and magnetic BEEM [49]. Therefore, we will focus here only on those aspects that are different, referring the reader to the literature [56, 67, 69–74] on the MTT for further information.

A first difference is that vacuum metal bonding, as used for the SVT, is not required. For the MTT, the base as well as the emitter structure can be created by vacuum deposition onto the collector semiconductor. For the fabrication of the thin tunnel barrier, methods well-established for fabricating magnetic tunnel junctions can be used [6]. Typically, this involves deposition of a 1–2 nm thin Al layer, and subsequent oxidation. The emitter is then deposited on top, and can be magnetically pinned using exchange bias from an antiferromagnet if desired.

A second notable difference is that in the MTT, the energy of the injected hot electrons is given by the voltage \( V_{EB} \) applied across the tunnel barrier. The hot-electron energy is thus tunable over a certain range of energies, typically up to 3 eV. This allows one to do spectroscopic studies. Also, as we have already seen in section 1.7, raising the hot-electron energy is beneficial for the transfer ratio of the device, since more electrons can be transmitted into the larger density of states higher up in the conduction band of the collector semiconductor. For the MTT, this is illustrated in the example of figure 24, which shows that above a certain threshold voltage set by the collector Schottky barrier, the transfer ratio grows more than linearly with \( V_{EB} \). The same behaviour is well-known in BEEM. Thus, in the MTT an increase of \( V_{EB} \) results in a larger collector current for two reasons, obviously the emitter current is increased, but at the same time the transfer ratio of the base is enhanced.

A third important aspect is that the MTT with a ferromagnetic emitter has only a single magnetic layer in the base, as opposed to two magnetic layers and the non-magnetic spacer for the SVT (and also for MTT with non-magnetic emitter). Since any base layer and interface is a source of scattering for the hot electrons, we anticipate that the transfer ratio of a MTT with magnetic emitter should be significantly higher. While the MTT was proposed as an alternative configuration, it is only after the understanding created by experiments with the SVT that it became clear that the MTT may allow larger output currents to be achieved. A drawback of the MTT with a ferromagnetic emitter is that the relative magnetic response (MC) is limited by the tunnelling spin polarization. This issue is discussed in the following paragraph.

The expressions for the transmission and magnetocurrent of the MTT with non-magnetic tunnel emitter (left design of figure 23) are the same as equations (2)–(7) for the SVT. However, the situation is different for the MTT with a ferromagnetic tunnel emitter, since the emitter current now becomes spin polarized, with the polarization determined by the tunnelling process. For the parallel magnetization state, the emitter tunnel current is a sum of a majority and a minority spin contribution according to [6]:

\[
I^P_E \propto \delta^M_E \delta^M_B + \delta^m_E \delta^m_B.
\]

Here \( \delta^M_E \) and \( \delta^m_E \) are the fraction of majority and minority tunnelling electrons associated with the emitter/insulator interface (label E), and \( \delta^M_B \) and \( \delta^m_B \) are the fraction of tunnel electrons associated with the insulator/base interface (label B). These are defined in the usual way with help of the tunnelling spin polarization \( P_t \):

\[
P_{t,E} = \frac{\delta^M_E - \delta^m_E}{\delta^M_E + \delta^m_E},
\]

and

\[
P_{t,B} = \frac{\delta^M_B - \delta^m_B}{\delta^M_B + \delta^m_B}.
\]
To obtain the collector current, we need to multiply each of the two emitter current components in equation (10) with the majority or minority spin hot-electron transmission factor appropriate for the ferromagnetic layer in the base (M2 in figure 23). We then have:

\[ I_C^P \propto \delta_M^E \delta_M^B T_M^E + \delta_m^E \delta_m^B T_m^E, \]

where \( T_M^E \) and \( T_m^E \) are defined as before in equations (4) and (5). If we assume that the emitter is magnetically pinned, then for the antiparallel case the magnetization of the base will be reversed and the resulting \( I_C^{AP} \) can be written as:

\[ I_C^{AP} \propto \delta_M^E \delta_m^B T_m^E + \delta_m^E \delta_M^B T_M^E. \]

Using these equations, the magnetocurrent can be cast into the familiar form:

\[ MC = \frac{I_C^P - I_C^{AP}}{I_C^{AP}} = \frac{2P_{E}P_{B}^*}{1 - P_{E}P_{B}^*}, \]

where we have defined a renormalized polarization \( P_{B}^* \) of the base as:

\[ P_{B}^* = \frac{\delta_M^E T_M^E - \delta_m^E T_m^E}{\delta_M^E T_M^E + \delta_m^E T_m^E}. \]

Note that this renormalized polarization is a combination of tunnelling factors and hot-electron transmission factors, and as such should not be confused with the regular tunnelling spin polarization as defined in equation (12). Thus, we see that the MC of a MTT with a ferromagnetic tunnel injector is determined by two completely unrelated physical processes. The polarization of the emitter current is due to spin-dependent tunnelling, while the subsequent transmission of the hot electrons in the base of the MTT is given by the spin-dependent scattering of hot electrons.

The fact that the MC depends on \( P_{E} \) has an important consequence, because the tunnel polarization \( P_{E} \) of the emitter/insulator interface is basically fixed, i.e. not dependent on the thickness of the emitter electrode. Rather, it is determined by the choice of metal and tunnel insulator [6]. In practice, this means \( P_{E} \) is limited to about 50% for typical ferromagnetic materials and \( \text{Al}_2\text{O}_3 \) tunnel barriers. Using the expressions given above it is straightforward to calculate the expected MC for a given emitter tunnel polarization (see figure 25, in which hot-electron transmission parameters for a Co base ferromagnet (\( \lambda_M^E = 21 \text{ Å} \) and \( \lambda_m^E = 8 \text{ Å} \) were used). As can be seen, the MC grows with increasing thickness of the base ferromagnetic layer as the base becomes a better spin 'analyzer'. However, at large base thickness essentially no minority spins are transmitted, and the MC saturates at a value that depends on the tunnel spin polarization. For the maximum tunnel polarization of 50%, the MC cannot be larger than 200%. This constitutes a fundamental difference with the SWT, where the MC can be made arbitrarily large by increasing the thickness of both the ferromagnetic layers of the base.

An interesting positive consequence is that the MTT can be used to probe the tunnel spin polarization at large bias. Recall that the MTT typically operates at a bias voltage of 1 V or larger. At such a voltage, the regular TRM of a ferromagnetic tunnel junction is usually negligibly small [6]. However, experimental MTTs with ferromagnetic emitter clearly display large magnetocurrent. An example is shown in figure 26 for a Co_{84}Fe_{16}/Al_{2}O_{3}/Co(10 nm)/Si MTT at 1 V bias and \( T = 77 \text{ K} \).

The emitter was pinned by Ir_{22}Mn_{78}. Top panel: the TRM of the emitter-base tunnel junction, measured at 10 mV and 1 V bias, respectively. Data by O M J van ‘t Erve, MESA’, in collaboration with S S P Parkin, IBM Almaden (unpublished).

![Figure 25. Calculated magnetic response of an MTT with ferromagnetic emitter, for different values of the tunnel spin polarization \( P_E \) of the emitter.](image-url)

![Figure 26. Bottom panel: collector current versus magnetic field for a Co_{84}Fe_{16}/Al_{2}O_{3}/Co(10 nm)/Si MTT at 1 V bias and \( T = 77 \text{ K} \). The emitter was pinned by Ir_{22}Mn_{78}. Top panel: the TRM of the emitter-base tunnel junction, measured at 10 mV and 1 V bias, respectively. Data by O M J van ‘t Erve, MESA’, in collaboration with S S P Parkin, IBM Almaden (unpublished).](image-url)
measured at a bias of 10 mV. This unambiguously proves that the disappearance of TMR at high bias does not necessarily imply that the tunnel current is no longer spin polarized.

3. Hot-electron spin injection

One of the crucial ingredients of semiconductor based spin electronics is the ability to inject a highly spin-polarized current into a semiconductor. While electrical spin injection appears conceptually simple, it has turned out to be a major experimental challenge. One successful approach is to use a paramagnetic II–VI semiconductor (Be$_{1-x}$Mn$_x$Zn$_{1-x}$Se) as spin aligner, yielding injection of highly polarized (90%) carriers [75]. The application potential of this method is however hampered by the large magnetic fields required to induce a sizeable spin splitting in the paramagnetic injector.

A second successful demonstration involved the injection from a ferromagnetic semiconductor Ga$_{1-x}$Mn$_x$As into GaAs [76]. Both approaches are limited to low temperature, although a room temperature magnetic semiconductor may change the prospects of this method. Nevertheless, the carrier spin polarization in any ferromagnet decays with temperature, especially when the Curie temperature is approached, such that the use of conventional ferromagnets with Curie temperature far exceeding room temperature is still the most promising route.

It has been argued [77] that direct spin injection from a ferromagnetic metal contact into a semiconductor, via diffusive transport, is nearly impossible due to the large mismatch of resistance between metal and semiconductor. This is in agreement with the absence of any sizeable magnetotransport effects in experiments ([77] and references therein). It was suggested [77–82] that the problem may be overcome either by using completely spin-polarized contacts (e.g. half-metallic ferromagnets), by using ballistic transport, or by introducing some kind of spin-dependent ‘interface resistance’. Such resistance can naturally exist at the metal/semiconductor interface due to differences in the electronic states, as shown for the case of Fe/InAs [83]. Alternatively, a tunnel barrier can be incorporated. The latter was invoked to explain the observed electrical injection of 2% spin-polarization from Fe into GaAs at room temperature [16]. More recently, tunnelling injection of spins across a Fe/AlGaAs Schottky contact was reported [17, 18], where a spin-polarization of 32% in the buried GaAs quantum well was obtained at low temperature. Spin injection across a thin AlO$_x$ tunnel barrier, from CoFe into an AlGaAs/GaAs semiconductor was also demonstrated [19, 20], the spin polarization being in excess of 11% at room temperature. Note that evidence for room temperature spin injection into GaAs across a vacuum tunnel barrier, from a Ni tip in a scanning tunnelling microscope, was reported in 1992 [84].

The work on hot-electron spin-transport in the SVT has led us to consider an alternative method to achieve controllable spin injection into semiconductors. The approach is based on spin-filtering of a current of hot electrons during transmission through a thin film of ferromagnetic metal, just prior to injection into the semiconductor. Since transport is based on non-equilibrium carriers, the conductivity mismatch does not play a role, and injection can be achieved with conventional ferromagnetic metals.

Figure 27 shows possible arrangements of the hot-electron spin filter (HESF) contacts. First, an unpolarized hot-electron current is generated using a Schottky barrier (left) or tunnel barrier (right). The hot-electron energy is determined by the Schottky barrier height (0.5–1 eV) or by the voltage (1–3 eV) applied across the tunnel barrier, respectively. Second, the hot-electron current is passed through a thin ferromagnetic metal film that acts as a polarizer for the hot-electron spins. The final step is the actual injection of the spin-polarized hot electrons into the semiconductor across a Schottky barrier. In this contact structure, the ferromagnetic film is generally (but not necessarily) separated from the semiconductor by a non-magnetic thin film, for example, to avoid magnetically dead layers or to tailor the Schottky barrier properties.

We emphasize that the conductivity mismatch [77] does not apply to hot-electron transport, as it is a non-equilibrium situation which cannot be described in terms of a series of resistances. Rather, when a current of hot electrons is injected into the ferromagnetic thin film, the energy and momentum relaxation of the hot electrons results in an exponential decay of the number of hot electrons with distance in the ferromagnet. Subsequent transmission into the semiconductor, across the Schottky barrier, occurs without any electric field or imbalance in electrochemical potential across the metal-semiconductor interface. In other words, the electron flow across the interface is due to the kinetic energy of the hot electrons, and generally no bias voltage is applied between the ferromagnet and the receiving semiconductor. This applies not only when transport is ballistic, but even if the hot-electron transport is diffusive, i.e. dominated by elastic scattering.

The two most important characteristics of the contact are the spin polarization of the injected electrons, and the fraction of the hot electrons that arrive in the semiconductor. We may assume that the spin polarization of the hot electrons is not affected by the transport across the non-magnetic metal film and its interface with the receiving semiconductor. This is reasonable since only non-magnetic materials are involved, and depolarization occurs via the weak spin-orbit scattering. It is then sufficient to determine the spin polarization after spin filtering by the ferromagnetic film, which has been done for a number of materials using the SVT and MTT. Note
that the large measured ratio between $I^P_C$ and $I^{AP}_C$ for the SVT: (i) shows that hot-electron spin filtering is extremely efficient, and (ii) proves that the spin polarization does not decay significantly during transport through the non-magnetic Au spacer between the Ni$_{80}$Fe$_{20}$ and Co layer.

Analysis of data such as shown in section 1.4 using equations (2)–(7) allows the determination of $\lambda^{M}_{\text{film}}$ and $\lambda^{m}_{\text{film}}$ from which the polarization of the hot electrons after spin-filtering can be evaluated. The resulting polarization as a function of film thickness for Ni$_{80}$Fe$_{20}$ or Co are shown in figure 28 (top panel). The spin polarization increases with increasing thickness, and rapidly approaches unity ($P = 96\%$ at 5 nm). Thus, almost complete spin-polarization is achieved for relatively thin films. The total transmission of a ferromagnetic film (bottom panel) is at an acceptable level, close to 0.1 for films of a few nanometres (the attenuation due to the interfaces was also taken into account).

For practical devices based on spin injection their properties at room temperature are relevant. Measurements such as those in section 1.6 at variable temperature show that hot-electron attenuation in ferromagnets is somewhat enhanced at room temperature due to scattering by thermal spin waves [51]. However, as can be seen in figure 28 (bottom panel), the net reduction of the transmission is relatively weak, especially for materials with a high Curie temperature such as Co. Besides the slight reduction of the transmission, thermal spin waves also produce mixing of the two spin channels [23]. This implies that the spin filter polarization also decays at higher temperature. Measurements of the MC versus temperature allow a direct quantification via equation (7). The data for a 6 nm Ni$_{80}$Fe$_{20}$ film in figure 29 shows that $P$ decays from 98% at low temperature, to 85% at room temperature. For Co, the decay of $P$ is negligible. Thus, even at room temperature, a HESF allows the creation of a highly spin-polarized current, with conventional ferromagnetic metals that can be manipulated with small magnetic fields. Several research groups are now working to directly measure the spin polarization of the hot electrons after injection into the semiconductor collector, using electroluminescence techniques [16–20, 75, 76]. This will show whether the assumption of negligible spin-depolarization of hot electrons at the metal/semiconductor interface is correct, and whether the high energy of the electrons gives rise to additional spin relaxation mechanisms in the semiconductor.

4. Outlook

Perhaps the most influential aspect of the SVT is that it showed the value of functional integration of semiconductor and ferromagnetic materials into hybrid electronic devices. Room temperature effects as large as 400% in small magnetic fields are now routinely obtained, a feature that has attracted significant attention. Much progress has been made in optimizing device performance, focusing on the output current level and noise sources. Further improvements are still required in order to capitalize on the huge magnetic sensitivity of such structures for application in magnetic field sensors or magnetic memories. In that context, the properties of SVT and MTT structures when scaled to sub-micron dimensions will be relevant.

At the same time, the SVT opened up a new route to systematically study the fundamental physics of spin-dependent transport of hot electrons at energies of the order of 1 eV. New insights into the origin of spin-dependent hot-electron scattering have been obtained, including the dominance of volume effects over interfaces in the spin-dependence of the transmission, the effect of thermal spin waves, and the surprisingly important role of elastic scattering processes. The latter feature points to the relevance of establishing the precise relation between scattering and structural properties of the devices, an area that has still remained largely unexplored. It also hints at an interesting route to the next generation of devices in which epitaxial growth techniques will be used to obtain transistors with high structural quality.

The basic concept of the SVT has also led to the development of a number of related device structures such as
the MIT. The tunability of the hot-electron energy provides unique possibilities for spectroscopic studies, while enhanced base transmission is anticipated. Moreover, it has been realized that hot-electron spin filtering may have some attractive features for spin-injection into a semiconductor, in particular the ability to reach a spin polarization near 100% with conventional ferromagnets. Work that is in progress will soon provide more information on the feasibility of this approach. Another interesting research avenue we are currently exploring is the study of hot-electron spin-transport in novel materials such as half-metallic ferromagnets and oxides. Such materials may also offer new types of hybrid electronic devices combining ferromagnets and semiconductors.

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