

Chapter 9

Spintronics and Nanomemory Systems

CNTs forest can be considered as a prototype of the magnetic memory, where ferromagnetic nanoparticles serve as cells of the magnetic memory – that is, ferromagnetic contacts are controlled by spin pulses, the transport of which is provided by nanotubes (see Fig. 9.1) [1]. This CNTs forest morphology can be obtained as a result of CVD growth using FM (ferromagnetic) nanodrops of predefined diameters and magnetism. In particular, we analyse Fe–Pt nanoparticle ability as nanomemory elements. Spin transport between magnetic cells is essential for nanomemory device operation.

There are similar spin valve structures formed from arrays of aligned carbon nanotubes. These devices require only one deposited ferromagnetic layer with the embedded iron catalyst nanoparticle serving as the other magnetic electrode [2]. Device magnetoresistance ratios are estimated about 25%. Spin valve structures include Au with Fe-catalyst drops, grown CNT array, Fe–Co layer, Cr–Au layer and glass substrate.

Spin-polarized transport over 5 μm long carbon nanotubes (CNTs) is observed in magnetically assembled spin valves with electrochemically deposited ferromagnetic (FM) electrodes. FM/CNT/FM devices as Co/CMT/Ni interfaces were considered. An annealing procedure is developed to produce stable devices with highly transmissive FM/CNT contacts [3].

9.1 Spin Transport Fundamentals

It has been found that carbon nanotubes can acquire magnetic properties when brought into contact with ferromagnetic materials. This property of carbon nanotubes has previously been predicted theoretically by M Ferreira and S Santivo. The effect is based on the exchange of polarized electrons between the ferromagnetic material and nanotubes. The main difficulty of the experiment was in detecting the weak magnetic moment of the nanotubes on the background of a

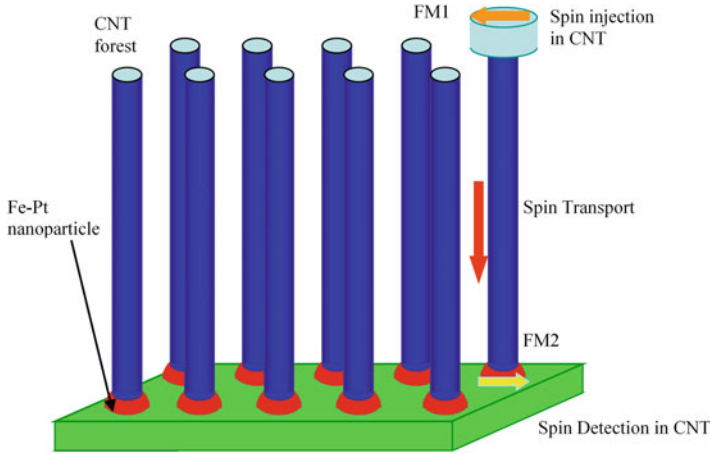


Fig. 9.1 CNT forest is grown on the Fe–Pt nanoparticles of the predefined radius as a fragment of magnetic nanomemory device, which can be realized if spin injection, spin transport and spin detection (spin recording) are provided

strong magnetic moment of the ferromagnetic sample. The sample used in the experiment was a thin film of cobalt or iron oxide and was uniformly magnetized strictly in one direction.

With the help of a magnetic force microscope, weak perturbing magnetic fields were revealed produced by nanotubes on the surface of the sample. According to the measurements, the magnetization of nanotubes is 0.1 Bohr magneton per one carbon atom, which is in contact with the film. The control experiments demonstrated that carbon nanotubes in contact with nonmagnetic materials such as Si or Au do not possess magnetism [4].

The general approach is to examine the chain of weakly interacting CNT spins in contact with the two ferromagnetic (FM) metals. Let us consider Hamiltonian $\mathcal{H}$ of the system of localized spins whose interaction is described by the model of Ising or Heisenberg:

$$\begin{aligned}\mathcal{H} &= -\frac{1}{2} \sum_{\mathbf{l}, \mathbf{l}'} J_{\mathbf{l}, \mathbf{l}'} [\alpha (S_{\mathbf{l}}^x S_{\mathbf{l}'}^x + S_{\mathbf{l}}^y S_{\mathbf{l}'}^y) + S_{\mathbf{l}}^z S_{\mathbf{l}'}^z] - \mu \sum_{\mathbf{l}} \mathbf{S}_{\mathbf{l}} \mathbf{H} = \\ &= -\frac{1}{2} \sum_{\mathbf{l}, \mathbf{l}'} J_{\mathbf{l}, \mathbf{l}'} [\alpha (S_{\mathbf{l}}^x S_{\mathbf{l}'}^x + S_{\mathbf{l}}^y S_{\mathbf{l}'}^y) + S_{\mathbf{l}}^z S_{\mathbf{l}'}^z] - \mu H \sum_{\mathbf{l}} S_{\mathbf{l}}^z,\end{aligned}\quad (9.1)$$

where $\mathbf{S}_{\mathbf{l}}$ is the spin localized in the lattice site $\mathbf{l}$, $\mathbf{H}$ is the external magnetic field, μ is the effective magneton and $J_{\mathbf{l}, \mathbf{l}'} = J(\mathbf{R}_{\mathbf{l}, \mathbf{l}'})$ is the energy of the exchange interaction (the exchange integral) of the nearest neighbours in the lattice. If $\alpha = 1$, Hamiltonian corresponds to the Heisenberg model, with $\alpha = 0$ – to the Ising model. The magnetic disordering in the system of CNT type is due to the weak magnetism of diamagnetic carbon atoms, where the only essential thing is the level

of magnetic interaction of the short-range order. In this case, a regular atomic structure of CNT is accompanied by a magnetic disorder of diamagnetic-liquids type. The correlation of the magnetic moments of CNTs can be ensured by a spin injection at the expense of the controlled magnetic field in the area of ferromagnetic contacts. At the same time, the initial values $J_{l,r}$ of the energy of the magnetic interaction till the controlled magnetic pulse have random values. The weak correlation between adjacent carbon atoms in the atomic structure of CNT is also supported by the excitation of FM–CMT–FM. Accordingly, in Eq. (9.1) summation is assumed only considering the efficiency degree of unbroken bonds.

In the Heisenberg model, the node approximation with the l -th node of the ideal crystal is associated with a localized magnetic moment proportional to the localized spin variable S_l .

In the case of paramagnetic materials, in the absence of an external magnetic field, S_l is oriented randomly from node to node and forms a system with magnetic disorder. Diamagnetic materials exhibit their magnetic properties due to the induced magnetic moments in the presence of a magnetic field. However, the magnetic moments are very small, and their correlated behaviour is characterized only by a short-range order. Therefore, in terms of spin transfer, paramagnetic and diamagnetic materials are quite similar.

In the case of cooperative spin interactions at low temperatures, magnetic order appears. It is characteristic for ferromagnets, antiferromagnets and ferrites. However, the order disappears at high temperatures and magnetic system itself is transformed into a paramagnetic with a certain level of magnetic disorder. Considering that S_l is a vector quantity, the types of disorder may be different. Of course, the thermal fluctuations significantly stimulate spin disordering.

Using the Ising model in its pure form (when $S_z = \pm 1$) for the systems such as FM–CMT–FM is not correct, because each node has three spin variables – S_x , S_y and S_z , which can take random values and are connected by the relation $S^2 = S_x^2 + S_y^2 + S_z^2$ and quantized by being inherently quasi-continuous. This is due to the boundary conditions on the FM–CNT interconnects. Really used constructive models of spin coupling (Ising model, classical and quantum Heisenberg model) for the system under investigation require further modifications. For example, in the near-field approximation, it is possible to use one parameter of the exchange interaction J in the model Hamiltonian:

$$\mathcal{H} = -\frac{1}{2} \sum_{l,l'} J S_l S_{l'} - \mu \mathbf{H} \sum_l \mathbf{S}_l. \quad (9.2)$$

Further steps may be associated with a version of Heisenberg magnetic system and the corresponding Hamiltonian:

$$\mathcal{H} = -\frac{1}{2} \sum_{\mathbf{l}, \mathbf{l}'} \xi_{\mathbf{l}} \xi_{\mathbf{l}'} J(\mathbf{R}_{\mathbf{l}\mathbf{l}'}) \mathbf{S}_{\mathbf{l}} \mathbf{S}_{\mathbf{l}'} - \mu \sum_{\mathbf{l}} \xi_{\mathbf{l}} \mathbf{S}_{\mathbf{l}} \mathbf{H}, \quad (9.3)$$

where $\xi_{\mathbf{l}}$ is the random value from 0 to 1, which can be adjusted by the temperature factor, that is $\xi_{\mathbf{l}} \sim \exp(-W/kT)$, where W is the magnetic interaction energy of the induced magnetic dipole with an external magnetic field.

Obviously, this type of complex Hamiltonian needs to be simplified. For example, we can neglect the spin deviations along the z axis and consider only the transverse deviations introducing spin variables such as $S^{\pm} = S_l^x \pm iS_l^y$.

The linearized equations of motion look as follows:

$$i\hbar \frac{dS_k^-}{dt} = 2SJ \sum_{k'} [S_k^-(t) - S_{k'}^-(t)], \quad (9.4)$$

where $S^2 = (S_l^x)^2 + (S_l^y)^2$.

Then, the approximate system of equations defining the dispersion of the model spin system will be as follows:

$$[\{2S \sum_k [J - \hbar\omega] S_k^- - 2S \sum_{k'} JS_{k'}^-] = 0. \quad (9.5)$$

This conceptual analysis allows us to imagine the complexity of the theoretical description of spin transport in the considered FM–CNT–FM systems and the amount of computational work.

9.2 Magnetoresistance Nanodevices

9.2.1 Spin Valve Concepts

Nanocarbon-based (e.g. CNT or graphene) spintronics is well-predicted theoretically. The main point of technological problem is the spin current injection efficiency into carbon nanotubes by using ferromagnetic electrodes. The next point of attention is spin transport experiments to obtain evaluations of digital information processing rates. A typical spintronic device is the so-called spin valve. The typical spin valve device contains two ferromagnetic electrodes (one ferromagnetic electrode for spin current injection and another for spin signal detection) and nonmagnetic material responsible for spin transport. For practical implementation, the spin signal should be large enough. A large spin signal can be obtained in a weak spin relaxation system, where spin can travel a long distance without being scattered in ballistic regimes. Nanocarbon-based materials (e.g. CMTs) due to long spin relaxation distances look as the best candidates in comparison with the nonmagnetic metals and semiconductors. The spin current is propagated for long

macroscopic distances due to some spin correlation. These nanocarbon-based materials are expected to be beneficial for implementation of large-scale spintronic logic circuits, quantum computing, fast data processing, magnetic hard disc drivers and dynamical nanomemory devices.

Electron transfer in spintronic devices depends on the spin polarization of the current. Such devices include a method of generating a spin-polarized current based on interconnects of nonmagnetic materials with ferromagnetic ones. In our case, these are Fe–Pt compositions.

EDOS spin splitting on the Fermi level defines the efficiency of spintronic devices. In accordance with Stoner model, the band structure of transition metals depends on spin orientation (up and down) [5]. The theoretical basis for considering spintronic devices is a magnetoresistance phenomenon. Namely, there are two variants of magnetoresistance – giant magnetoresistance (GMR) and tunnelling magnetoresistance (TMR).

Electronic transport through CNTs depends on the contacts with electrodes. For short CNTs at low temperatures, we should also pay attention to Kondo effect – the formation of a many-body dynamical singlet between a localized spin and delocalized conduction electrons of electrodes [6]. It should be borne in mind that carbon nanotubes possess in addition to spin also orbital degeneracy. The increase of degeneracy corresponds to the enhancement of Kondo temperature, which is important for potential applications.

Giant magnetoresistance is a quantum mechanical effect observed in thin metal films consisting of alternating ferromagnetic and nonmagnetic conductive layers. The effect is a substantial change in the electrical resistance of such a structure when changing the relative direction of magnetization of adjacent magnetic layers. The direction of magnetization can be controlled, for example, by applying an external magnetic field. The effect is based on the scattering of electrons depending on the spin direction.

The GMR was discovered in 1988 (2007 Nobel Prize, Albert Fert and Peter Grünberg) as a large change in resistance of magnetic Fe/Cr multilayer in the presence of an applied magnetic field [7]. Soon after GMR was discovered in Fe/Cr/Fe trilayers [8]. As it was shown later, the effect can be obtained in trilayers having other magnetic materials such as Co [9].

Such trilayer structures, i.e. sandwiches of two ferromagnetic metals separated by a thin spacer layer of normal metal (Fig. 9.2), are of great industrial importance. They are called spin valves and are used as magnetic field sensors. The resistance of the device is dependent on the relative magnetization orientation of the ferromagnets. It is R_p when the magnetizations are parallel and R_A when they are antiparallel. The GMR ratio is defined as $GMR = \frac{R_A - R_p}{R_p}$.

The TMR is a quantum mechanical effect manifests itself when a current flows between the two ferromagnetic layers separated by a thin (about 1 nm) dielectric layer. The total resistance of the device, the current which flows due to the tunnelling effect, depends on the relative orientation of the two fields, the magnetization of the magnetic layers and resistance above at the perpendicular

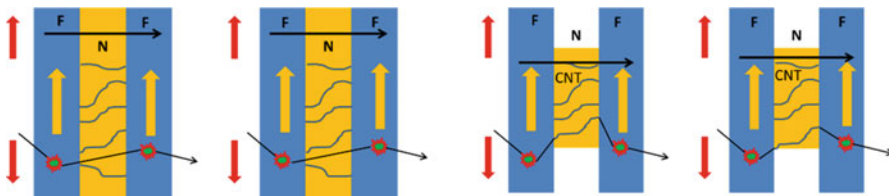


Fig. 9.2 Giant magnetoresistance (GMR) device. A thin normal metal spacer (N) separates two ferromagnets (F), the current flows perpendicular to the plane of the sample and N space can be filled by introduced (or grown) CNT, e.g. metal-like

magnetization layer. The effect of the tunnel magnetic resistance is similar to the effect of a giant magnetic resistance, except for the use of insulating layer of the tunnel barrier instead of a nonmagnetic metal layer.

The TMR was discovered in 1975 by Jullière [10] in a device that consisted of two Fe films separated by Ge. It was first in 1995 when room temperature TMR was discovered by Moodera [11] and caused a great interest in spintronics after GMR was discovered in 1988 [6]. Our idea is that the same effect can be reached by introducing into the N space semiconductor CNTs.

The TMR signal operates in the same way as the GMR $TMR = \frac{R_A - R_P}{R_P}$, where $R_P = (1/G_P)$ and $R_A = (1/G_A)$ are the resistances (conductances) of the device for parallel and antiparallel orientations, respectively, of the ferromagnet magnetization. However, the nature of TMR is different than that of the GMR.

The resistance for the antiparallel magnetization is normally higher than for the parallel one. At the applied zero magnetic field, the relative orientation of the magnetization is governed by the exchange coupling between the ferromagnetic layers. The sign and size of the interlayer exchange coupling is dependent on the thickness of the nonmagnetic spacer. The coupling can thus be ferromagnetic and antiferromagnetic dependent on the spacer thickness [12, 13].

9.2.2 Spintronic Device Descriptions

The first two-terminal CNT spin valve device was presented in [14], where spin-dependent transport was demonstrated for MWCNTs with 9% magnetoresistance ratio (MR) at 4.2 K.

The device shown in Fig. 9.2 is the so-called current perpendicular to the plane (CPP) geometry. The resistance of such geometry is very low and difficult to detect. For practical applications, structures with the current in the plane (CIP) are used because they have higher resistance and thus higher difference with the magnetic field [9].

The GMR can be understood through the Mott's two-current model [15]. According to that model, the electrical conductivity of metal can be described by two more or less independent channels, one for majority spins and the other for

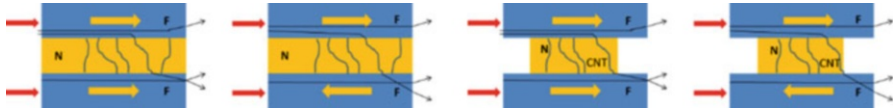


Fig. 9.3 Current in plane (CIP) spin valve and equivalent resistor mode, separate channels are for minority and majority spins, the electrons scatter from one F layer to the other on the way through the sandwich

minority spins. Scattering processes that conserve spin states are much more probable than the processes that flip spins. Another view proposed by Mott is that the scattering probability of spin up and spin down is quite different, independent on the nature of the scattering process [16]. This is shown schematically in Fig. 9.2.

The difference in the resistance of ferromagnets can be explained by the exchange split band structure. The scattering of the electrons depends on where the electron band crosses the Fermi level [9].

As mentioned above, commercial spin valves have normally the current in plane (CIP) geometry. This gives higher resistance and thus higher resistance difference between parallel and antiparallel spin orientations. Having the current in plane of the sandwich will qualitatively give the same effect as in CPP devices. This situation is shown schematically in Fig. 9.3. Electrons with minority and majority spins are treated separately. When the electrons flow through the sandwich, they will scatter back and forth from the upper F layer to the lower one.

We consider a possibility for the creation of magnetic memory devices on the magnetoresistance phenomenon (GMR or TMR) consisting of sandwiches of two ferromagnetic metals separated by a thin spacer layer of *normal metal* (see, e.g. Fig. 9.3) or *semiconductor*. Spin valves are of great industrial importance and are used as magnetic field sensors. The resistance of the device is dependent on the relative magnetization orientation of the ferromagnets. The TMR signal operates in the same way as the GMR, where R_P and R_A are the resistances of the device for parallel and antiparallel orientations, respectively, of the ferromagnets magnetization (Figs. 9.2 and 9.3) [17]. The main idea is to reach the same effect by introducing metal or semiconductor-like CNTs into the N space after magnetically controlled CVD CNTs growth [1, 17].

9.3 Magnetic Disorder and Spin Transport

We outline two main problems of spintronics, namely, spin states of localized Fe–Pt ferromagnetic spaces and spin transport efficiency provided by CNTs system.

9.3.1 Magnetic Disorder in Fe–Pt Nanodrops

Consider ferromagnets with localized magnetic moments in which the magnetic atoms are randomly distributed at the sites of the crystal lattice. It is also advisable to consider structurally non-regular metal alloys such as metal glasses A_xB_{1-x} , on which magnetic (spin) disorder is superimposed. This can affect the additional contributions to the exchange integrals of the spin interaction. Such ferromagnetic systems are technologically important, in particular for the CVD magnetically controlled CNTs growth, when nanodrops of ferromagnetic binary metallic glass are used. This is essential also for creating prototypes of magnetic nanomemory systems. The consideration of two possible opposite situations for magnetic properties A_xB_{1-x} metal glasses is essential and should be carried out.

The First Model Only the nearest lattice neighbours interact. In this case, the magnetic order exists only if the concentration of magnetic atoms x exceeds the critical value x_c (percolation threshold), i.e. in addition to the temperature phase transition, a concentration phase transition at $x > x_c$ occurs at zero temperature.

The Second Model The exchange interaction is non-zero for any distance between magnetic atoms. In this case, ferromagnetism exists for any arbitrarily small concentration x . We will assume that the exchange potential decreases with distance exponentially $V(r) \sim \exp(-r/R)$, where r is the average distance between magnetic atoms and R is the radius of the exchange potential.

The magnetism of such systems can strongly differ from the properties of ordinary ordered ferromagnets. The behaviour of such model spin systems at various concentrations x is of interest.

9.3.1.1 Diluted Ferromagnets with the Nearest Neighbour Interaction

Consider a system of localized spins, the interaction of which is described by the Ising or Heisenberg model (see Eq. 9.1). The disorder in this system is most easily created in two ways. First, it is possible randomly replace magnetic atoms in the lattice sites by nonmagnetic atoms. In this case, $J_{ij} = V_0 p_i p_j$, where $p_i = 1$, if the node is occupied by a magnetic atom, and $p_i = 0$, if it is nonmagnetic. Secondly, it is possible to break sites in a random way, i.e. consider some J_{ij} to be zero. Such a situation can in principle be realized in ferromagnets with indirect exchange through nonmagnetic ions. Removing the latter or replacing them with ions that do not take part in indirect exchange corresponds to a rupture of the connection between the spins.

At $T = 0$ K, all spins, bound by the exchange interaction, are oriented in parallel. If the concentration x of magnetic atoms or unbroken bonds is small, the magnetic atoms form clusters isolated from each other. The dimensions of them do not change with increasing of the crystal volume. We call such clusters finite.

The average magnetic moment of the whole crystal is zero in this case. With increasing x , the characteristic size of finite clusters increases, and at a certain critical concentration, x_c arises a cluster, in which the number of particles is proportional to the volume of the system. We shall henceforth call such a cluster infinite.

The magnetic moment of the crystal $M(x)$ is different from zero at $x > x_c$ and equals $M(x) = M(1)P(x)$, where $M(1)$ is the moment of the ordered crystal and $P(x)$ is the concentration of magnetic atoms (per lattice site) included in an infinite cluster. Thus, in a disordered ferromagnet at $T = 0$ K, a phase transition occurs from the paramagnetic state to the ferromagnetic state when the concentration $x = x_c$. The problem of the concentration phase transition, or, the problem of the origin of an infinite cluster, is studied in percolation theory (see, e.g. [18, 19]). In accordance with the above two methods of creating disorder in the lattice, it is customary to distinguish between the *site problem* and the *node problem*. The critical concentration is called the *percolation threshold*. The values depend on the type of lattice and are different for the *site problem* and the *node problem*.

They were obtained as a result of computer modelling. For cubic lattices in the *site problem model* $x_c = 0.243$ (*body-centred cubic*-bcc) and 0.195 (*face-centred cubic*-fcc); in the *node problem model* $x_c = 0.178$ (bcc) and 0.120 (fcc). The detailed table containing x_c for different lattices is given in the review of Essam [19]. The value $P(x)$ changes from 0 at $x = x_c$ to 1 at $x = 1$. Results of typical simulation of behaviour are presented in Fig. 9.4.

The next problem is to determine the concentration dependence of the *rigidity* of the magnons D . It is convenient to use the coupling that exists between the concentration dependences and the conductivity of an equivalent resistance lattice (see Kirkpatrick [18]): $\frac{D(x)}{D(1)} = \frac{\sigma(x)}{P(x)\sigma(1)}$.

The equivalent resistance network is constructed from a magnetic lattice in the following way: the magnetic bonds V_{ij} correspond to conductivities distributed according to the same law as σ_{ij} . According to Kirchhoff's laws, the time dependence of the potential $V_i(t)$ at the i -th node is determined by the following equation:

$$\frac{dV_i(t)}{dt} = -\left(\frac{\sigma_0}{C}\right) \sum_{i'} \{V_i(t) - V_{i'}(t)\}, \quad (9.6)$$

where C is the magnet assembly capacity and σ_0 is the conductivity of each connecting link. This equation is similar to the equation for magnon modes:

$$i\hbar \frac{dS_i^-(t)}{dt} = 2SJ \sum_{i'} \{S_i^-(t) - S_{i'}^-(t)\}. \quad (9.7)$$

Thus, the solution of the conductivity problem with respect to a random resistance network can be related to the properties of the eigenfunctions $|\alpha\rangle$ of the residual matrix:

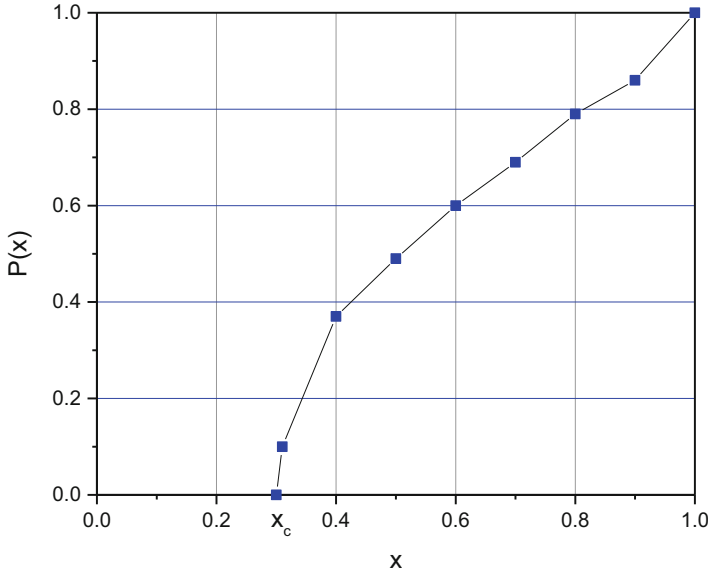


Fig. 9.4 Typical behaviour of $P(x)$ simulation for simple cubic lattice

$$\Delta = \sum_{i,i'} \{|i\rangle\langle i|- \} \{ |i\rangle\langle i'| \}. \quad (9.8)$$

The macroscopic static conductivity $\sigma(p)$ of the network should be equal the limiting value of the dissipative part of the response function to the external field for zero frequency and infinitely long waves. The classical analogue of the fluctuation-dissipative theorem directly leads us to an expression equivalent to the Kubo–Greenwood formula [20–23]:

$$\sigma(x) = \lim_{\omega \rightarrow +0} \left\{ \lim_{q \rightarrow 0} \left[\left(\frac{\omega^2 C^2}{q^2} \right) \sum_{\alpha} |\langle \mathbf{q} | \alpha \rangle|^2 \frac{\lambda_{\alpha}}{\lambda_{\alpha}^2 + \omega^2 C^2} \right] \right\}. \quad (9.9)$$

In the limiting case of long waves, the Fourier transform $\langle \mathbf{q} | \alpha \rangle$ of the eigenfunction of the matrix Δ reduces to the Fourier transform of the plane wave corresponding to the momentum $\mathbf{q}_{\alpha} = \langle \alpha; \mathbf{R} | (\hbar/i) \nabla | \mathbf{R}; \alpha \rangle$, and $\mathbf{R}$ is the position of the i -th node. In the continuum approximation (see [18]), the eigenvalues of the matrix are estimated as $\lambda_{\alpha} \approx a^2 q_{\alpha}^2 D(x)$, where $D(x)$ characterizes the concentration of the magnetic nodes of the system and a is the lattice constant.

Using the relationship $\lambda_{\alpha} \approx a^2 q_{\alpha}^2 D(p)$ can be found [18]: $\sigma(x) = \sigma_0 a^2 D(x) P(x)$.

It expresses the relationship between the rigidity parameter of spin waves in a dilute ferromagnet $D(x)$ and the volume conductivity of the corresponding network

$\sigma(x)$. Mobility percolation $\mu(x) = \sigma(x)/P(x)$ is not constant but decreases to zero at $p \rightarrow p_c$ (see [24]).

Note that the magnetic ordering in a dilute Heisenberg ferromagnet becomes unstable with respect to the excitation of long-wave magnons exactly at the percolation threshold p_c . This confirms the well-founded hypothesis that under these conditions, the value really determines the point of the magnetic transition in the system at $T_C \rightarrow 0$ [25].

9.3.1.2 Bethe Lattice of the Spin System

The system of spins randomly located in the Bethe lattice is an exactly solvable model of a disordered ferromagnet. The Bethe lattice is a convenient topological model of atomic structures of a certain connectivity. In this case, the lattice sites and connections can be associated with certain local physical characteristics.

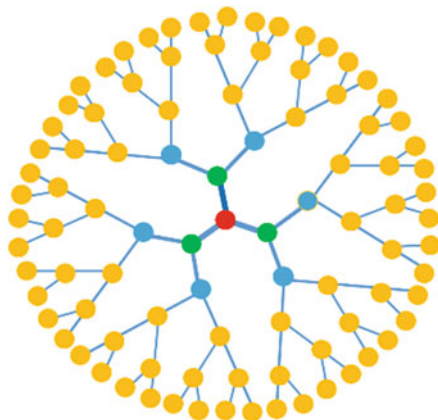
The Bethe lattice is a convenient topological model of atomic structures of a certain connectivity. In this case, the lattice sites and connections can be associated with certain local physical characteristics (see Fig. 9.5). In the Bethe lattice, z sites emerge from each node, giving rise to disjoint branches (see Fig. 9.5). Suppose, that at the sites of the Bethe lattice, the magnetic atoms are randomly distributed, whose concentration (calculated on the lattice site) is equal x and whose interaction is described by the Hamiltonian (9.1).

Let us find the expression for the mean magnetic moment of the lattice at $T \ll V_0$. Let R be the probability that a branch originating from some fixed node (*node 0* in Fig. 9.5) is finite, i.e. the magnetic atoms in it do not form an infinite cluster. The branch can be finite if the node adjacent to the null (*node 0*) in the given branch (*node 1*) is not occupied (probability $(1 - x)$) or if this node is busy, but the outgoing branches $z - 1$ are finite (probability). Consequently, R satisfies Eq. [28–33]: $R = 1 - x + xR^{z-1}$. The simplicity of this equation is due to the main property of the Bethe lattice is a non-intersection of branches. The average moment is $M(x) \sim P(x) = 1 - R^z$. The equation $R = 1 - x + xR^{z-1}$ always has a trivial solution $R = 1$ corresponding to $M = 0$. A nontrivial solution $R \neq 1$ appears at $R \neq 1$, x_c is the percolation threshold in the Bethe lattice. Near the threshold $(x - x_c)/x_c \ll 1$, from we have:

$$\frac{M(x)}{M(1)} = \frac{2z}{z-2} \frac{x - x_c}{x_c}. \quad (9.10)$$

Stinchcombe was the first who studied spin waves in the Bethe lattice [26, 27]. He used the relationship $\frac{D(x)}{D(1)} = \frac{\sigma(x)}{P(x)\sigma(1)}$. He determined the conductivity by fixing the potential difference between the beginning (point 0, Fig. 9.5, Bethe tree) and the boundary. It was shown that for the Bethe lattice the effective medium method is valid with respect to the z^{-1} parameter everywhere, except for the region near the threshold, for which the conductivity is proportional $(x - x_c)^2$. de Gennes

Fig. 9.5 Bethe lattice with $z = 3$



paid attention to the fact that macroscopic conductivity should be determined for a fixed external field in the entire sample [30]. In this case, near the threshold, the conductivity varies according to the law [28–30]: $\sigma \sim (x - x_c)^3$ (the Bethe lattice with the index $z = 3$). It then follows that $D \sim (x - x_c)^2$. We see that the exact results for the Bethe lattice are completely consistent with the similarity hypothesis near the percolation threshold.

9.3.1.3 Spin Waves

Crangle and Boszort et al. [34, 35, 37] discovered an astonishing phenomenon: the paramagnetic Pd metal becomes ferromagnetic with little Fe or Co doping. Later it was shown that the spontaneous moment arises even at concentrations x of Fe and Co of the order of 10^{-4} [36], as well as in alloys $\text{Pd}_{1-x}\text{Mn}_x$ [38].

The properties of the investigated magnetic system are completely determined by the ratio of two characteristic lengths: the radius R of the potential and the average distance between the impurities $r_c \approx n^{-1/3}$, where n is the number of spins per unit volume. If $R \gg r_c$, then a large number of magnetic atoms $\nu_R = (4/3)\pi n R^3 \gg 1$ effectively interact with this spin. Therefore, the energy of the exchange interaction of each spin with the environment depends little on the specific configuration. In this case, naturally, is proportional to the average energy of the exchange interaction, i.e. $T_C \sim x$ [39]. A completely different situation arises when $R \ll r_c$, i.e. $\nu_R \ll 1$. Now the energy of the exchange interaction of a given spin with the remaining spins depends very strongly on the specific configuration. In fact, for most spins, the nearest neighbour is located at a distance of order r_c . In this case, even insignificant fluctuations in the arrangement of the spins lead to a significant change in the energy of the exchange interaction [40]. In a system of randomly located points, all connections between points that are separated by a distance $r < r_{\max}$ are included. When $r_{\max} \ll n^{-1/3}$, there are only finite clusters of connected atoms in the system. For almost all atoms are connected in an infinite cluster. At

what value $r_{\max}$ does the infinite cluster first appear? According to [41], the critical value $r_{\max}$ is equal to $r_0 = (0,87 \pm 0,01)n^{-1/3}$. At high temperatures $r(T) < r_0$ and the system is paramagnetic. The temperature of the ferromagnetic transition is determined from the condition $r(T) = r_0$. Then, $T_C \approx V_0 S^2 \exp\left(-\frac{0,87}{Rn^{1/3}}\right) = V_C \exp\left(-\frac{0,87}{Rn^{1/3}}\right)$, where S is the spin of the system, $r(T) = R \ln \frac{V_0}{T}$ and $V(r) = V_0 \exp(-r/R)$ [40].

Rigidity coefficient of spin waves. To determine the concentration dependence of the rigidity D for small concentrations of magnetic atoms, it is convenient to use the coupling between D and the conductivity σ of an equivalent resistance network. The generalization of the formula $\frac{D(x)}{D(1)} = \frac{\sigma(x)}{P(x)\sigma(1)}$ to a disordered ferromagnet in which the spin interaction potential depends on the distance r according to the law $V_{\text{ferro}}(r) = U_0 \frac{R}{r} \exp(-r/R)$ consists in the fact that the probability $P(x)$ should be replaced by the number of magnetic atoms whose binding energy exceeds the energy of the spin wave. Consequently, in the low-frequency limit [40] $D \sim \frac{\sigma}{n}$.

Here σ can be regarded as the conductivity of electrons in a system of chaotically located centres, in which the probability of electron hopping between centres is determined by the formula $V_{\text{ferro}} = U_0 \frac{R}{r} \exp(-r/R)$, where $U_0 = J^2 N(\epsilon_F) \frac{\Omega_0}{2\pi a^2 R}$ [40]. Its dependence on concentration is known: $\sigma \sim r_0^{-(2+\nu)} \exp\left(\frac{0,87}{Rn^{1/3}}\right)$ [42],

where ν is the index of the correlation radius and $r_0 = 0,87n^{-1/3}$. Then follows $D \sim n^{-(1-\nu)/3} \exp\left(\frac{0,87}{Rn^{1/3}}\right)$ [40]. It should be noted D/T_C that the concentration dependence is determined only by the index of the correlation radius ν : $(D/T_C) \sim n^{-(2-\nu)/3}$, where $\nu \approx 0.70 - 0.85$ (for $\text{Pt}_{1-x}\text{Fe}_x$). Assuming $n \sim (x - x_c)$ for small concentrations, we obtain $\frac{T_C}{D} \sim x^{(2-\nu)/3}$. Further, we get $\ln\left(\frac{T_C}{D}\right) \approx A \frac{(2-\nu)}{3} x$ where A is the proportionality coefficient (see Fig. 9.6).

Magnetic atoms are randomly distributed at the lattice sites. Their concentration is low ($x \ll 1$). Their exchange interaction is described by the Heisenberg Hamiltonian.

Spin Transport Mechanisms There are mainly two kinds of geometries for CNTs spin valve devices: local geometry and non-local geometry. The local geometry has two ferromagnetic electrodes; the spin current is injected from one electrode, then transported through CNTs and detected by another electrode.

The spin current is mixed with the charge current. The spin signal is detected as the difference in resistance between the parallel state and antiparallel state of two magnetic electrodes. Generally, it can be expressed as magnetoresistance (MR) as $MR = (R_A - R_P)/R_P$. Usually, the two magnetic electrodes should be different in geometric aspect ratios to obtain different coercive fields; thus the two electrodes can be aligned to antiparallel or parallel states by sweeping the magnetic field.

Tsukagoshi et al. fabricated the first two-terminal CNT spin valve device. Spin-dependent transport was demonstrated through MWNTs, and $MR = 0.09$ at 4.2 K was observed [14].

Spin waves transport can be tested experimentally using corresponding schemes.

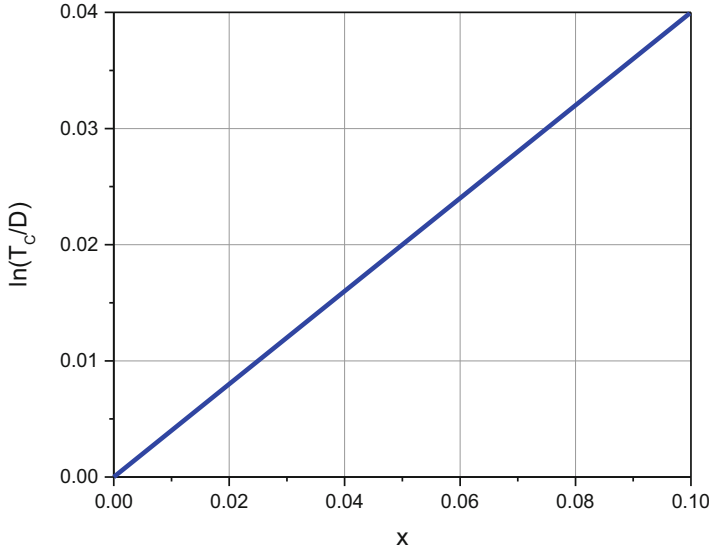


Fig. 9.6 Typical model dependence $\ln(T_C/D)$ via the concentration x alloys for $\text{Pt}_{1-x}\text{Fe}_x$ ($\nu = 0.8$, $x \ll 1$)

The simplest measurement scheme is a local spin valve (see Fig. 9.7a). However, the non-local scheme is more correct (see Fig. 9.7b, c). The first variant of the schematic diagram of non-local spin valve device FM1, FM2, FM3 and FM4 denotes the four electrodes, respectively. Current I is injected into CNT from FM3 and extracted from electrode FM4. The voltage is detected between electrode FM1 and FM2 (see Fig. 9.7b). The second variant of schematic diagram of non-local spin valve device FM1, FM2, FM3 and FM4 denotes the four electrodes, respectively. Current I is injected into CNT from FM1 and extracted from electrode FM4. The voltage is detected between electrode FM2 and FM3 (see Fig. 9.7c). We also can consider end-type FM–CNT contacts taking into account a set of end-type schemes in Fig. 9.8, which corresponds to spin transport in the case of CNT forest (see Fig. 9.1).

9.3.1.4 Spin Injection and Detection

Consider the mechanism of electrical spin injection and detection in a nonmagnetic material. We pay attention to traditional conventional two probe spin valve and the non-local four probe techniques.

Electrical Spin Injection and Detection A current which moves through a diffusive conductor can be considered to be carried in parallel by two independent spin channels, a channel for the spin-up electrons and one for the spin-down electrons [43]. This holds when most scattering events preserve the spin orientation;

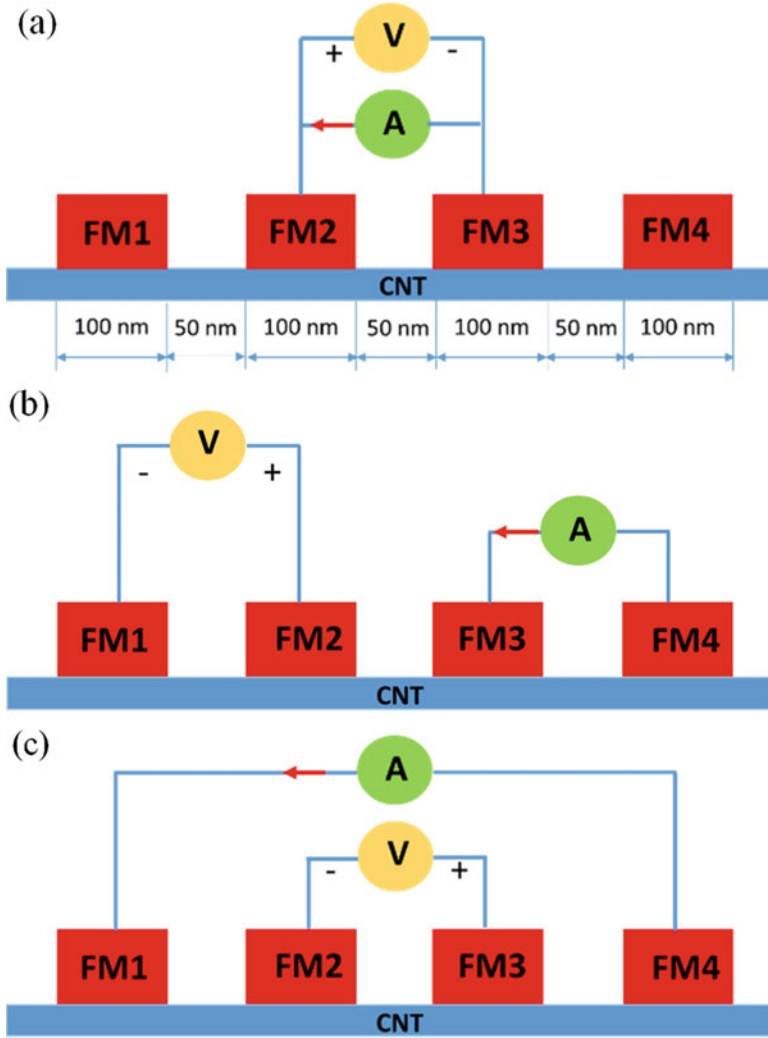


Fig. 9.7 (a) Local side-type CNT spin valve device; (b) non-local side-type CNT spin valve device, the first scheme; (c) non-local side-type CNT spin valve device, the second scheme

otherwise, the two channels become strongly coupled and are not independent anymore. The spin current density $j_{\uparrow\downarrow}$ of the spin-up $\uparrow$ or spin-down $\downarrow$ channel is proportional to the gradient of the electrochemical potential of one of the two spin species and is given by the relation $j_{\uparrow\downarrow} = \frac{\sigma_{\uparrow\downarrow} \partial \mu_{\uparrow\downarrow}}{e \partial x}$, where $\sigma_{\uparrow\downarrow}$ is the conductivity, $\mu_{\uparrow\downarrow}$ is the chemical potential of the two spin channels and e is the electron charge. The total charge current is given by $j = j_{\uparrow} + j_{\downarrow}$, and the spin current j_s is given by $j_s = j_{\uparrow} - j_{\downarrow}$. Looking now into a ferromagnet in which $\sigma_{\uparrow} \neq \sigma_{\downarrow}$, we can define a spin polarization P given by the relation $P = \frac{j_{\uparrow} - j_{\downarrow}}{j_{\uparrow} + j_{\downarrow}} = \frac{\sigma_{\uparrow} - \sigma_{\downarrow}}{\sigma_{\uparrow} + \sigma_{\downarrow}}$; using $\sigma = \sigma_{\uparrow} + \sigma_{\downarrow}$ which

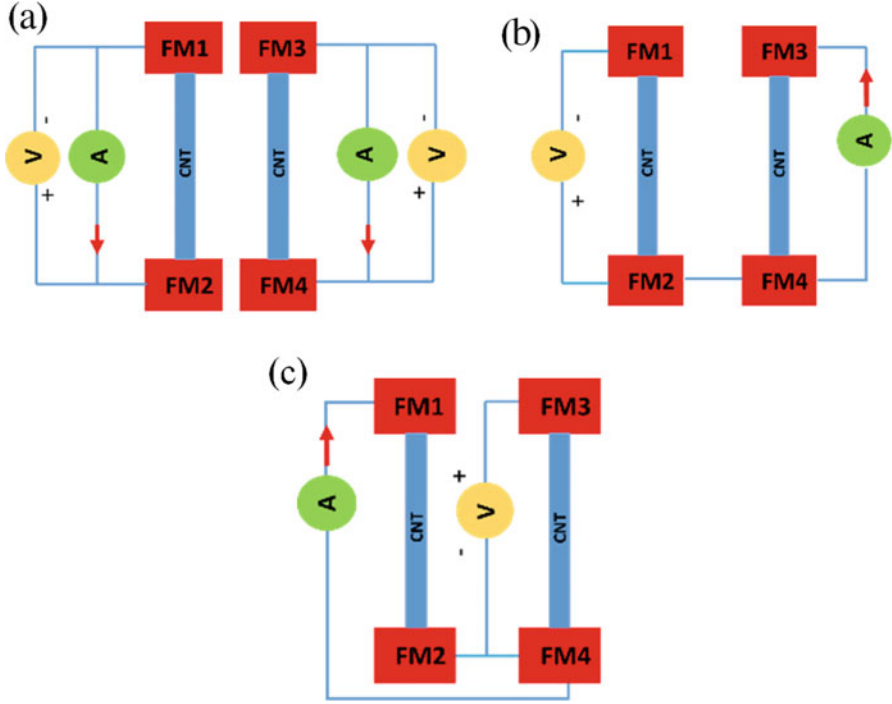


Fig. 9.8 (a) Local end-type CNT spin valve device; (b) non-local end-type CNT spin valve device, the first scheme; (c) non-local end-type CNT spin valve device, the second scheme

is the total conductivity in the system we obtain $j_{\uparrow} = (1 + P) \frac{\sigma \partial \mu_{\uparrow}}{e \partial x}$ and $j_{\downarrow} = (1 + P) \frac{\sigma \partial \mu_{\downarrow}}{e \partial x}$.

We can use the ferromagnet as a spin source, because $j_{\uparrow} \neq j_{\downarrow}$. In practice, ferromagnets are used in conventional two-terminal spin valve device. A spin valve device contains two ferromagnetic electrodes, the injector and the detector, which make contact to a nonmagnetic material. Sending an electrical current through the ferromagnetic injector creates a spin accumulation in the nonmagnetic material. In a diffusive material, this accumulation decays exponentially with distance from the spin injector following the differential equation $\frac{\partial^2 \Delta \mu}{\partial x^2} = \frac{\Delta \mu}{\lambda_s^2}$,

where $\lambda_s = \sqrt{D_s \tau_s}$ is the spin relaxation length, D_s is the diffusion constant, τ_s is the spin relaxation time and $\Delta \mu = \mu_{\uparrow} - \mu_{\downarrow}$. If the spins relax fast, such as to reach a uniform distribution over spin direction before arriving at the spin injector, then the spin detector measures the average chemical potential which is zero. In this case, switching the magnetization of the ferromagnets from parallel to antiparallel gives no change in the total resistance of the spin valve. However, in the situation in which the detector is set within a distance λ_s from the spin injector, a switching of the magnetization of the ferromagnets from parallel to antiparallel produces a

change in the spin valve resistance. The reason for this is that when the magnetization of injector and detector are parallel, the spins can enter easily the spin detector, corresponding to a low resistance of the spin valve. However, when the magnetization of the detector is antiparallel to the injector, then the spins first have to flip their magnetization before they can enter the detector; this results to an increase in resistance.

Spin Injection and Detection, the Non-local Technique The conventional two-terminal spin valve geometry can be used to extract the spin relaxation length λ_s . However, this technique is not so reliable in respect of pure spin transport detection. Really we can observe mixed signals (e.g. Hall effects, anisotropic magnetoresistance interference effects, magneto-coulomb effects). A four-terminal non-local spin valve geometry [44, 45] is able to completely separate the spin current path from the charge current path. Hence, the signal measured is due to spin transport only. To determine spin accumulation in the non-local geometry, one needs to contact the system under investigation (metal, semiconductor, carbon nanotube) with four electrodes. At least two of these should be ferromagnetic. They act as spin injector and spin detector, respectively [44, 45].

As shown in Fig. 9.7b, two of the ferromagnetic electrodes are connected to a current demonstrates the spin transport in a four-terminal CNT-based spin valve device using the non-local spin valve geometry: a current I is injected from electrode FM3 through the contact barrier into CNT and is extracted at contact FM4. The voltage difference is measured between contact FM2 and FM1. The non-local resistance is $R_{\text{non-local}} = (V_+ - V_-)/I$. Injection of up spins by contact FM3 results in an accumulation of spin-up electrons underneath contact FM3, with a corresponding deficit of spin-down electrons. Due to spin relaxation the spin density decays on a scale given by the spin relaxation length. The electric voltage is measured by contact FM1 and FM2 in the ideal case of 100% spin selectivity. A positive non-local resistance is measured. In cases of spin injection and spin diffusion for antiparallel magnetizations the voltage contacts probe in opposite spin directions gives a negative non-local resistance.

9.3.1.5 Spin Precession

In the non-local geometry but also in the conventional geometry, the effectiveness of spin injection and detection strongly depends on the relative magnitude of the contact resistance of the ferromagnetic electrodes compared to the resistance of the nonmagnetic material. A small contact resistance with respect to the nonmagnetic material, unavoidably, results in the so-called conductivity mismatch and also provides an extra path for spin relaxation at the ferromagnetic contacts [46]. This has to be taken into account in fitting of the spin precession measurements by solving the one-dimensional Bloch Eqs. [45]. The Bloch equations describe the

combined effect of diffusion, precession and spin relaxation in the system $D_s \frac{d^2 \vec{\mu}}{dx^2} - \frac{\vec{\mu}}{\tau} + \frac{g\mu_B}{\hbar} (\vec{B} \times \vec{\mu}) = 0$, where D_s is the diffusion constant, $\vec{\mu}$ is the chemical potential vector of the spin species, τ is the spin relaxation time, g is the g -factor, $\vec{B}$ is the magnetic field, $\hbar$ is Planck's constant and μ_B is the Bohr magneton.

9.3.1.6 Spin Relaxation

Consider possible spin relaxation mechanisms in a nonmagnetic material, for example, in a single wall nanotube. Spin relaxation in a nonmagnetic material is an unavoidable process which brings a non-equilibrium population of spins (found at the spin injector interface to the nonmagnetic system) into a uniform distribution over spin directions [47]. This description is also adequate for corresponding graphene-based systems. Four important mechanisms of spin relaxation in semiconductors should be analysed (see, Fig. 9.9, [47]).

Elliott–Yafet Mechanism Conduction electron spins can relax via ordinary momentum scattering from impurities, boundaries and phonons if the lattice ions induce spin-orbit coupling in the system ([47], Fig. 9.9a). Spin-orbit interaction mixes spin-up and spin-down states, and therefore the Bloch states (momentum eigenstates) are not eigenstates anymore. The result is that the longitudinal time, T_1 , is proportional to the momentum scattering time, τ_p . Therefore, the spin flip length $\lambda_s = \sqrt{D_s T_1}$ is proportional to the mean free, since the diffusion constant is proportional to τ_p .

Dyakonov–Perel Mechanism Spin dephasing occurs because electrons feel an effective magnetic field, resulting from the lack of inversion symmetry and from the spin-orbit interaction, which changes in random directions every time the electron scatters to a different momentum state (Fig. 9.9b, [47]).

Bir, Aronov and Pikus Mechanism The electron–hole exchange scattering can lead to efficient electron spin relaxation in p-type semiconductors [48, 49]. According to the electron–hole exchange interaction, within the elastic scattering approximation, the spin lifetime limited by the Bir–Aronov–Pikus mechanism is given by the Fermi Golden Rule [50].

Hyperfine Interaction Mechanism The hyperfine interaction is the interaction between the magnetic moments of the electrons and the nuclear magnetic moment (Fig. 9.9d, [47]). This mechanism dominates in the case of localized electrons, for example, in quantum wells and quantum dots.

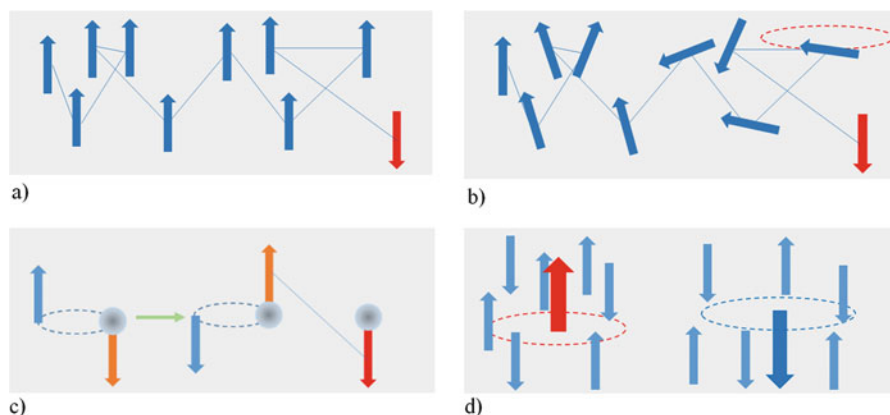


Fig. 9.9 Four important mechanisms of spin relaxation in semiconductors. (a) Elliott–Yafet mechanism. (b) Dyakonov–Perel mechanism. (c) Bir–Aronov–Pikus mechanism. (d) Hyperfine interaction

9.4 Concluding Remarks

CNTs are expected to be good candidates for spintronic devices. The first organic spintronic device was reported by Tsukagoshi et al. in 1999 [51], and remarkably, it consisted of a multiwalled carbon nanotube contacted by Co contacts. Many other spintronic experiments on multi- and single-walled carbon nanotubes followed up. Unfortunately, the experiment of Tsukagoshi and all other experiments performed after his work have made use of the conventional two-terminal spin valve geometry. Magnetoresistance phenomenon (GMR and TMR) for nanomemory devices based on CNTs of various morphologies (i.e. various chiralities, diameters) including metal- and semiconductor-like ones and ferromagnetic Fe–Pt contacts can be potentially viewed as an alternative for electromagnetic nanosensing and magnetic nanomemory.

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