38

Spin-Transfer Nano-Oscillators

38.1 Introduction

John Slonczewski (1996) and Luc Berger (1996) predicted that electron currents in magnetic multilayer devices can transport angular momentum from one magnetic layer to another, thereby exerting a torque on the local magnetization (see Figure 38.1a). Under the correct conditions, they predicted that the magnetization can undergo sustained oscillations at microwave frequencies. The necessary conditions include (1) the spacer layers between the magnetic layers are thin, <50 nm, so that spins do not depolarize as they go from one layer to another; (2) the device is sufficiently small (<100 nm) so that the amount of spin momentum transported by the electron current is a significant fraction of the angular momentum of the magnetic element; (3) there are sufficient nonlinearities in the configuration to stabilize the precessional orbits. Microwave oscillations in magnetic multilayers were first detected by Tsoi et al. (1998, 2000) in a point-contact geometry and then microwave emission was directly observed in nanoscale device structures by Kiselev et al. (2003) and Rippard et al. (2004). These structures are commonly referred to as spin-torque oscillators, spin-transfer oscillators, or spin-transfer nano-oscillators (STNOs). The most common device configuration, shown in Figure 38.1a, uses two conducting magnetic layers: a polarizer whose magnetization is fixed and a free layer whose magnetization is free to rotate in response to the current generated spin torque. The layers are separated by a metallic or an insulating spacer layer. An example of microwave emission from an STNO for various bias currents is shown in Figure 38.1b.

The advantages of spin-transfer oscillators are that they are highly tunable by current and magnetic field, they are among the smallest microwave oscillators yet developed, they are relatively easy to fabricate in large quantities, they are compatible with standard silicon processing, and they operate over a broad temperature range. STNOs are closely related to giant magnetoresistance (GMR) and tunneling magnetoresistance (TMR) devices that have been developed for magnetic recording read heads and magnetic random access memory. STNOs utilize both spin-transfer torque, to set the magnetization into oscillation, and the GMR or TMR effect, to produce an output voltage. Challenges still remain before widespread applications of STNOs are possible. These challenges include increasing the output power above the present value of ~0.5 μW (Deac et al. 2008), removing the need for large applied magnetic fields, understanding and controlling the oscillator linewidth, and reducing device-to-device variations.

This article presents a tutorial on spin-transfer oscillators; no prior knowledge of spin transfer or magnetic devices is assumed. For more in-depth discussion of spin-transfer effects the reader is referred to reviews by Sun (2006), Silva and Rippard (2008), and Ralph and Stiles (2008).
38.2 Physics of Giant Magnetoresistance, Tunneling Magnetoresistance, and Spin Transfer

The physics of spin-transfer torques is similar to that giving rise to the GMR and TMR effects. These effects arise from the fact that, in addition to charge, electrons have an additional degree of freedom, the electron spin. In ferromagnets (and sometimes in normal metals and semiconductors), the valence electrons can be polarized: there are more spins pointing in one direction rather than another. The spin density polarization along a particular axis $\hat{z}$ is given by $P = (n_\uparrow - n_\downarrow)/(n_\uparrow + n_\downarrow)$, where $n_\uparrow$, $n_\downarrow$ are the number of valence electrons with their spins aligned and anti-aligned along $\hat{z}$. The properties of the conduction electrons, including their momentum state and scattering rates, can depend on the spin direction. Since an electron has a charge $e$, the electron spin $\vec{S}$ is associated with a magnetic moment $\vec{\mu} = g\gamma\vec{S}$, where $g = (ge/2m_e)$ is the gyromagnetic ratio, $m_e$ is the electron mass, and $g$ is the gyromagnetic constant, which for the materials considered here, is close to the free electron value of $g = 2$. In general, the magnetic moment has both an orbital as well as a spin component. For the materials currently used in STNOs, the orbital moment is not important and we consider just the moment due to spin. The charge on an electron is negative so $\gamma$ is negative and $\vec{S}$ and $\vec{\mu}$ are antiparallel. Since the electron has a spin of $s = \hbar/2$ the electron moment is nearly equal to the Bohr magneton $\mu_B = (ge/2m_e)(\hbar/2) = \mu_B = (\hbar c/2m_e) = 9.274 \times 10^{-23} \text{ A m}^2$. For an electron in a metal the gyromagnetic ratio is approximately (it will vary depending on the material)

\[\gamma = 1.76 \times 10^{11} \text{ A m}^2/\text{J s}.\]

Often it is more convenient to remember the value $\gamma/2\pi = 28 \text{ GHz/T}$ since this will be the constant relating the electron spin precession and STNO oscillator frequency to the effective magnetic field. To characterize magnetic devices it is useful to evaluate the local magnetization $\vec{M}$ which is the moment per unit volume: $\vec{M} = \mu_0 \sum \mu$, where $V$ is the volume of the element being considered and $\mu$ indexes over all of the moments in $V$. Often, as in Figure 38.1a, we refer to the normalized moment of a device layer using lower case symbols $\vec{m} = \vec{M}/|\vec{M}|$, where $\vec{m}$ is a unit vector in the direction of the magnetization.

Polarization of electrons is most prominent in magnetic materials, which have a spontaneous magnetic moment: there is a net alignment of the electron spins on each atom and a well-defined ordering of moments from atom to atom. The alignment of the electron spins is due to a quantum mechanical effect that requires the electron wave function to be anti-symmetric. The anti-symmetry of the wave function leads to an effective interaction, the exchange interaction, that lowers the system energy if the electron spins on a given atom, if given a choice, are aligned. This energy reduction is due to the requirement that the wave function for the electrons must go to zero if two electrons with the same spin occupy the same location. The vanishing of the probability that two electrons with the same spin state occupy the same location reduces the Coulomb repulsion energy. For electrons on the same atom, this effect leads to the alignment of spins in partially filled $d$ and $f$ orbitals as described by Hund’s rules. For electrons on different atoms, there are similar effects, however, the exchange interaction can be either positive (favoring spin alignment) or negative (favoring spin anti-alignment), depending on the relative distance between the atoms and their relative orientation. If the inter-atomic exchange interaction favors alignment then the

---

* In this article all equations and quantities are expressed in Système International (SI) units.
material, below a Curie temperature $T_C$, is a ferromagnet and if the inter-atomic exchange interaction favors anti-alignment then the material, below a Néel temperature $T_N$, is an antiferromagnet.

For STNO, GMR, and TMR devices we are most interested in conducting ferromagnets that have both a spontaneous polarization as well as mobile electrons. The simplest model of a conducting ferromagnetic is the $sd$ model proposed by Mott (1936). In this model, we consider transition metals to consist of partially filled $d$ shells, which contain most of the polarization, and mobile $s$ electrons that carry most of the current as illustrated in Figure 38.2. The $d$ electrons are relatively localized; they have large effective masses and low velocities. The $d$ electrons are polarized since the shell is partially filled and the exchange interaction promotes alignment of the electron spins. For Fe, Co, Ni metals (which are the most common materials used in the magnetic layers shown in Figure 38.1a), the moment per atom at 0 K is of 2.22, 1.72, 0.61 $\mu_B$, respectively, corresponding to roughly 0.5–2 more electron spins pointing parallel to the magnetization axis than antiparallel. The saturation magnetization $M_s$ and Curie temperatures $T_C$ for Fe, Co, Ni metals at 0 K are $1.71 \times 10^6$, $1.42 \times 10^6$, $0.48 \times 10^6$ A/m, and 770°C, 1131°C, and 358°C, respectively.

Calculated spin-dependent band density of states and conductivities for face centered cubic (fcc) Co are shown in Figure 38.2b. Most of the polarization is due to $d$-like electrons well below the Fermi surface. Electrons with magnetic moments parallel or antiparallel to the magnetization direction are referred to as majority or minority electrons, respectively. Note that for Co at the Fermi surface there are more minority spins than majority spins since the minority $d$ bands are at the Fermi surface. However, the conductivity at the Fermi surface is dominated by the majority $s$-like electrons. This can be intuitively understood by realizing that most of the electron-scattering events in clean metals preserve the spin direction, i.e., scattering events that cause spin flips are relatively rare. Hence, majority electrons must scatter into other majority states and similarly minority electrons must scatter into other minority states. Since there is a larger density of states at the minority Fermi surface there are more states to scatter into for this channel. This excess of scattering, referred to as $sd$ scattering, causes the conductance in the minority state $s$-band to be considerably smaller than for electrons in the majority state.

In a rigid band picture, the states for the minority electrons are similar to that for the majority electrons except they are shifted by additional exchange energy. At the Fermi surface, the minority electrons will have smaller wave vectors and velocities. For metallic contacts, this potential barrier, as seen in Figure 38.3, causes scattering at the interface between regions with different magnetizations. In particular, since there is a large potential step for spins that go from being majority to minority, and vice versa, these electrons will be scattered strongly. The strong scattering leads to a higher resistance when the magnetizations of the free and polarizer layers are not aligned (Camley and Barnaś 1989, Grünberg 2007). The device resistance is then a function of the relative angle $\theta_m$ between the magnetizations in the two layers and, to a good approximation, is given by $R = R_{av} - (\Delta R/2)\cos(\theta_m)$, as shown in Figure 38.3d, where $R_{av} = (R_{sp} + R_p)/2$, $\Delta R = R_{sp} - R_p$ and $R_{sp}$, $R_p$.

**FIGURE 38.2** (a) Schematic density of states for a transition metal. (b) Calculated density of states for majority and minority electrons in fcc Co. (c) Conductivity of majority electrons. (d) Conductivity of minority electrons. At the Fermi surface, most of the current is carried by majority $sp$ electrons. (From Tsymbal, E.Yu. and Pettifor, D.G., *Phys. Rev. B*, 54, 15314, 1996. With permission.)
are the device resistances in the antiparallel and parallel states, respectively. The geometry shown in Figure 38.3 is referred to as current perpendicular to the plane (CPP) configuration. GMR originally was discovered in structures that had the current flowing in the plane (CIP) of the layers. Spin-transfer devices require large spin-polarized currents flowing from one layer to another and, hence, must be in the CPP configuration. The resistance for a metallic device with a single ferromagnetic interface is typically a few ohms and the resistance change is a few tenths of an ohm. The fractional resistance change, the magnetoresistance (MR) ratio, is \( \frac{\Delta R}{R} = \frac{R_{\text{ap}} - R_p}{R_p} \approx 2\% - 20\% \). In practice, it is impossible to obtain a sharp change in the magnetization if there is a strong exchange coupling at the interface between the two magnetic layers so a thin spacer layer is inserted. The states in the spacer layer must match that of the conducting electrodes or there will be unwanted scattering due to the spacer layer. The most common choice of spacer layer in metallic devices is Cu. The spacer layer thickness must be less than the spin-flip relaxation length, which for copper at room temperature is approximately \( \lambda_{sf} \approx 100-400 \text{ nm} \) (Albert et al. 2002).

An alternative to inserting a thin metallic spacer layer is to insert an insulator that is thin enough to allow tunneling through the barrier. The resulting device is referred to as a magnetic tunnel junction (MTJ). The device resistance is now determined by both the tunneling process (how the wavefunctions decay in the barrier) and by the density of available states in the electrodes. If the wave function decay is independent of whether the electron spin is majority or minority the MR ratio is given by Julliere’s formula (Julliere 1975), which accounts only for density of states effects:

\[
\frac{\Delta R}{R} = \frac{2P_1P_2}{1 - P_1P_2}
\]

where \( P_1, P_2 \) are the polarizations of the electrons at the Fermi surface of the two layers that are contributing to the tunneling current. For typical ferromagnetic metals with polarizations of 20\%–40\%, this will give MR ratio values on the order of 10\%–40\%.

If the polarizations could be made close to 100\%, as in a half metal which has only one spin polarization at the Fermi surface, then the MR ratio could get arbitrarily large. However, the MR ratio, as it is conventionally defined, is not a good figure of merit for an STNO. As discussed in Section 38.6, the maximum output power of an STNO is \( P_{\text{out}} \propto I^2 R L / \Delta R^2 \), where \( R_L \) is the load resistance, and hence \( \Delta R / R_{\text{av}} \), which varies between 0 and 2, is a better figure of merit.
The TMR can be greatly enhanced if the majority and minority wavefunctions decay differently in the tunnel barrier (Butler et al. 2001, Yuasa and Djayaprawira 2007). This is not due to any explicit dependence of tunneling on the electron spin but rather due to the fact that majority and minority electrons have, at the Fermi surface, different spatial wavefunctions with different symmetries and different decay lengths. To exploit this effect, there has to be a high degree of crystalline texture in the electrodes and the tunnel barrier and the barrier has to be thick enough to allow for the wave function of one channel, usually the minority spin channel, to decay in magnitude substantially below that of the other channel. The most successful MTJ s to date have used (001) MgO barriers with recrystallized body centered cubic (bcc) CoFeB electrodes, which have MR ratios of 200%–500% at room temperature. Typical magnetoresistive response, MR ratios, and barrier resistivities are shown in Figure 38.4. For STNO devices, thin barriers are required with resistance area products RA < 10 \( \Omega \mu \text{m}^2 \). As seen in Figure 38.4b, the MR ratio falls off when barriers are this thin and the tunnel barriers are very susceptible to breakdown, making the fabrication of MTJs for STNO applications a challenge.

The discussion above describes why the device resistance will depend on the relative orientation of the magnetizations in the two magnetic electrodes. If the local magnetization acts on the flow of the conduction electrons then, conversely, the flow of conduction electrons must act on the local magnetization. This is shown schematically in Figure 38.3c where the polarizer and free layer magnetizations are at an angle \( \theta_m \) relative to each other. Conduction electrons are polarized by flowing through the polarization layer then, when they are incident on the free layer, their polarization rotates to align with the direction of the free-layer magnetization. This rotation of the conduction electron moment is due to a torque applied by the local magnetization and, hence, the local magnetization feels an equal and opposite torque. The transverse component of the momentum of the conduction electron is transferred to the local magnetization. The spin-transfer torque on the free layer, which is the change in spin momentum of the localized moment per second, is given by two terms, a transverse torque (which is shown in Figure 38.3c) and a field-like torque. The transverse torque is the most important and is given by

\[
\bar{\tau}_i = \frac{g_s(\theta_m) \mu_e I}{ye} \vec{m}_f \times (\vec{m}_f \times \vec{m}_p)
\]

where

- \( I/e \) is the number of electrons incident per second on the free layer
- \( \mu_e/\gamma = \hbar/2 \) is the magnitude of the spin of each electron
- \( g_s(\theta_m) \) is the spin-torque parameter which is proportional to the polarization of the spin currents
- \( \vec{m}_f \times (\vec{m}_f \times \vec{m}_p) \) is a vector with magnitude \( \sin(\theta_m) \) and direction perpendicular to \( \vec{m}_f \) in the \( \vec{m}_f, \vec{m}_p \) plane

It will be important to keep track of the sign of the applied current. Here, a positive current will mean that electrons are

---

**FIGURE 38.4** Magnetoresistance of MgO MTJs from Yuasa et al. (2004). (a) Resistance vs. applied field showing a high resistance antiparallel state and a low resistance parallel state. (b) The MR as a function of MgO barrier thickness. (c) The RA product vs. MgO barrier thickness. (From Yuasa, S. et al., Nat. Mater., 3, 868, 2004. With permission.)
flowing from the free layer into the polarizer and a negative current will mean that electrons are flowing from the polarizer into the free layer. For this convention, a negative current means that the free-layer magnetic moment is pushed toward that of the polarizer, as the case in Figure 38.3c; a positive current means that the magnetic moment is pushed away from that of the polarizer.

While we have been focusing on the torque applied to the free layer by electrons flowing from the polarizer, there is also a torque on the polarizer caused by the flow of electrons. This can be seen in Figure 38.3c, which shows that the electrons that are reflected from the free layer form a polarized current incident on the polarizer, which will apply a torque to that layer. The current induced torques will drive the magnetizations in each layer in the same direction: if the current is negative, as shown in Figure 38.3c, the free layer magnetization will be pushed toward the polarizer and the polarizer will be pushed away from the free layer magnetization. In fact, the only thing that distinguishes the polarizer from the free layer is that, in device construction, we try to make the polarizer magnetization fixed and unresponsive to the current induced torques. Some STNO designs utilize the fact that all the layers will have a torque to set multiple layers into motion (Tsoi et al. 2004).

A complete theory of spin-dependent transport must encompass both the spin torque and GMR/TMR. A useful theory by Slonczewski (2002) predicts the magnitude of the spin torque for symmetric metallic devices in terms of two dimensionless parameters $P_{\text{sc}}$, $\Lambda$:

$$g_s(\theta) = \frac{P_{\text{sc}} \Lambda^2}{(\Lambda^2 + 1) + (\Lambda^2 - 1)\cos(\theta_m)}$$

$$\Lambda^2 = \frac{G R_p \mu_B}{R_p} \mu_B$$

where $G = A e^2 k_B / 4\pi^2 \hbar$ is a conductance with a magnitude of $\sim 5 S$ for a 75 nm device with a copper spacer layer, $A$ is the area of the device and $k_B$ is the Fermi vector of the spacer layer. $P_{\text{sc}}$ is the spin current polarization (the difference between the currents in the two of the spin channels divided by the total current), which is proportional to $\delta R = R_{\text{up}} - R_{\text{down}}$. Given $R_{\text{up}}$ and $R_{\text{down}}$, one can calculate $\Lambda$ and $P_{\text{sc}}$ and, hence, the spin torque. This model explicitly demonstrates that spin torque, device resistance, and magnetoresistance are intimately related. Typical values for $\Lambda$ and $P_{\text{sc}}$ for the materials currently used in STNOs are 1.5–4 and 0.2–0.6, respectively. Plots of $g_s$, $g$, $\sin(\theta_m)$, and $R$ are shown in Figure 38.3d for a standard all-metallic device. Xiao et al. (2004) have shown that the above expression for $g_s(\theta)$ is in reasonable agreement with more quantitative calculations using realistic band structures and Boltzmann transport theory and have extended the theory to include asymmetric structures. A few things to note: there is no spin torque when the layers are aligned since $\sin(\theta_m)$ goes to zero; $g_s$ is typically larger near 180° when the polarizer and free layer are close to being anti-aligned. The asymmetry in the spin torque can be understood by realizing that the polarized electrons need not be transmitted to create a torque. When the layers are close to anti-aligned, the torque can be very large since most of the electrons are reflected and there are many more spin-transfer events per transmitted electron.

### 38.3 Single-Domain Equation of Motion and Phase Diagrams

#### 38.3.1 Equation of Motion without Damping

A localized spin, $i$, obeys the torque equation, which says that the rate of change of angular momentum $\vec{S}_i$ is equal to the torque applied $\vec{\tau}$:

$$\frac{d\vec{S}_i}{dt} = \vec{\tau}_i$$

In the case of a conservative system, a system that does not exchange energy with its environment, we can write the torque in terms of an effective field $H_{\text{eff}}$ acting on the $i$th spin

$$\frac{d\vec{S}_i}{dt} = -\mu_0 \vec{\gamma} \left( \vec{M} \times H_{\text{eff}} \right)$$

Since $\vec{\mu} = \gamma \vec{S}_i$, the equation can be rewritten in terms of the magnetic moment

$$\frac{d\vec{M}_i}{dt} = -\mu_0 \vec{\gamma} \left( \vec{M} \times H_{\text{eff}} \right)$$

where $\vec{M}(\vec{r})$ is the magnetization or the moment per unit volume at a point $\vec{r}$. The effective field is the negative gradient of the free energy density with respect to the magnetization*:

$$H_{\text{eff}} = -\nabla \mu U = -\nabla \mu \left( U_{\text{ap}} + U_{\text{an}} + U_{\text{ex}} + \cdots \right)$$

$$= \vec{H}_{\text{ap}} + \vec{H}_{\text{an}} + \vec{H}_{\text{ex}} + \cdots$$

For STNO devices the most important energy terms are the interaction with the applied field (Zeeman energy) $U_{\text{ap}} = -\mu_0 \vec{M} \cdot \vec{H}_{\text{ap}}$, the magnetostrictive energy due to dipolar interactions between the spins $U_{\text{an}}$, the anisotropy energy due to crystalline or interfacial energies $U_{\text{ex}}$, and the exchange energy due to spin-dependent quantum mechanical interactions $U_{\text{ex}}$. The effective field is then the sum of the applied magnetic field $H_{\text{ap}}$, magnetostrictive energy, Zeeman energy, and anisotropy energy.

* Two fields that are of interest in magnetism: the magnetic field $\vec{H}$ in units of A/m and the magnetic flux density $\vec{B}$ in units of tesla. $\vec{H}$ and $\vec{B}$ are related through $\vec{B} = \mu_0 \vec{H} + \mu_0 \vec{M}$. It is convenient to think of $\vec{H}$ as the “driving” field and $\vec{B}$ as the actual local field a small test spin (such as a proton) would feel. Extending this concept, we view $H_{\text{eff}}$ as an effective driving field, not the local flux density. We often quote the magnetic field as $\mu_0 \vec{H}$ in tesla since this is a more convenient and easily understood unit. It is useful to remember the correspondence: 1 T $\Rightarrow$ 796 kA/m. A good introduction to magnetostrictic energy, demagnetizing factors, anisotropy energies, and exchange energies can be found in O’Handley 2000.
field $\vec{H}_{an}$, anisotropy field $\vec{H}_{an}$, and exchange field $\vec{H}_{ex}$. In this section, we will assume that the exchange interaction is sufficiently strong that all of the moments in a device layer are aligned and move together. This approximation is often referred to as the single domain or macrospin model. In this model the magnitude of the moment is constant in time. We will discuss the case of nonuniform moments in Section 38.8. We will also assume, for the time being, that the polarization layer moment is truly fixed and only consider the dynamics in the free layer. We write the equation of motion in terms of the normalized moment of the free layer $\vec{m}_f = \vec{M}_f / M_s$:

$$\frac{d\vec{m}_f}{dt} = -\mu_0 \gamma \vec{m}_f \times \vec{H}_{eff}$$

The equation says simply that the free layer moment will precess, with constant magnitude, around the effective field always in a direction that is perpendicular to the energy gradient and hence on a constant energy trajectory.

The magnetostatic energy density and field for a uniformly magnetized element can be written as $U_{ms} = \frac{1}{2} \mu_0 M_s^2 \left( N_{xx} m_x^2 + N_{yy} m_y^2 + N_{zz} m_z^2 \right)$ where $N_{xx}, N_{yy}, N_{zz}$ are the shape-dependent demagnetizing factors along the $x, y, z$ directions, respectively. The demagnetizing factors sum to 1 and are smallest for directions along the longest dimensions of the magnetic element. This energy term simply states that the moment likes to lie along the longest dimension of the element. For a sphere the demagnetizing factors are $N_{xx} = N_{yy} = N_{zz} = 1/3$ and there is no preferred direction that minimizes the magnetostatic energy. For an infinitely long cylinder, the demagnetizing factors are 1/2 perpendicular to the axis of the cylinder and 0 along the axis so that the moment likes to lie along the axis of the cylinder. For a typical $75 \times 50 \times 3$ nm STNO free layer, shown in Figure 38.5, the demagnetizing factors are $N_{xx} = 0.047, N_{yy} = 0.072$, $N_{zz} = 0.881$ and the moment likes to lie in the plane of the film along the long axis of the device. The magnetostatic field is obtained by differentiating the energy with respect to the magnetization: $\vec{H}_{ms} = -M_s \left( N_{xx} m_x + N_{yy} m_y + N_{zz} m_z \right)$. The magnetostatic field along a particular axis is proportional to the magnetization along that axis and points opposite to the direction of the magnetization.

The magnetostatic energy difference between the $x$ and $y$ directions at $T = 300$ K for the device shown in Figure 38.5, normalized to $k_B T$, is $\Delta U_{ms} / k_B T = 27$. It is a general feature of these nanoscale devices that the characteristic energies are not much greater than the thermal energies and hence thermal fluctuation are an important part of device operation.

The anisotropy energy may be due either to the preference for the moment to lie along certain crystalline directions or for the moment to lie perpendicular to interfaces. An example of the former is in L10 FePt films in which the magnetization likes to lie along the $c$-axis of the face-centered tetragonal cell. An example of the latter are Co/Ni multilayers in which the magnetization likes to be perpendicular to the thin film interfaces.

![Figure 38.5](image-url) Free precessional orbits of the free layer of a magnetic device with a $3 \times 50 \times 75$ nm, $M_s = 800$ kA/m free layer (a) with no damping and no applied field, (b) with no damping and 1.25 T applied field in the $z$-direction, (c) with no damping and 0.3 T applied field 45° off the $z$-axis, and (d) with damping ($\alpha = 0.02$) and no applied field.
Here, we consider only uniaxial anisotropy energies of the form $U_{an} = -\frac{1}{2} \mu_0 M H_{an} (\mathbf{m} \cdot \hat{z})^2$ where $\pm \hat{z}$ are the easy axis directions (the low energy directions assuming $H_{an}$ is positive). The anisotropy field is given by $H_{an} = M_m \hat{z}$. Typical anisotropy energies and anisotropy fields are $U_{an} = 0.5 - 1 \times 10^6 / \text{m}^3$ and $\mu_0 H_{an} = 1.25 - 2.5 \text{T}$. The form of the uniaxial anisotropy energy term is similar to that for the $z$-component of the magnetostatic energy.

If the magnetization is oriented along a particular direction and allowed to evolve according to the torque equation, the magnetization will precess about the effective fields and trace out a constant energy orbit. Some typical "free precession" orbits are shown in Figure 38.5 for a magnetic device without perpendicular anisotropy and in Figure 38.6 for a magnetic device with perpendicular anisotropy. Also shown are the precession frequencies. These "free precession" orbits will be approximately those sampled by the spin-transfer driven oscillations discussed later. In Figure 38.5a, there are no applied fields and the moment precesses around the internal magnetostatic and anisotropy fields, while in Figure 38.5b and c the moment precesses around a combination of internal fields and applied fields. For the case of weak fields and no perpendicular anisotropy (Figure 38.5a), there are two different types of orbits: in-plane orbits and out-of-plane orbits. In addition, for high symmetry geometries there can be degenerate orbits, i.e., two orbits with the same shape and frequency. For a circular orbit, the frequency is given by the precession speed divided by the orbit circumference $f_0 \equiv \frac{\gamma \mu_0}{2 \pi m_i} |\mathbf{m} \times H_{eff}|$ where $m_i$ is the orbit radius. For an elliptical orbit, the precession frequency is proportional to the geometric mean of the precession speed at the major and minor axes of the orbit divided by the orbit circumference

$$f_i \equiv \frac{\gamma \mu_0}{2 \pi \sqrt{m_a m_b}} |\mathbf{m} \times H_{eff}| \frac{m_a \times H_{eff}}{m_b \times H_{eff}} ,$$

where $m_a$ and $m_b$ are the major and minor axis lengths, respectively. For the perpendicular orbits shown in Figure 38.5a and b, which are nearly circular, the resonant frequencies are given by

$$f_i \equiv \frac{\gamma \mu_0}{2 \pi} (H_{ap} - N_x M_x) = \frac{\gamma \mu_0}{2 \pi} (H_{ap} - N_x M_x \cos(\theta)) ,$$

where $\theta$ is the angle between the moment and the $z$ axis and we have assumed $H_{ap} = H_{ap} = H_{ap}$ and $N_x, N_y$ are close to zero. For zero applied field, the frequency is negative denoting clockwise rotation (as viewed from the positive $z$ direction) while, for large applied fields the frequency is positive denoting counterclockwise rotation. For zero applied field, the frequency decreases as $\theta$ increases and the magnetostatic field decreases. Near $\theta = 90^\circ$ the orbit frequencies get small and the orbits go in-plane. For large $z$-axis applied fields, $H_{ap} > M_s$, the frequency will increase as $\theta$ increases going to

![FIGURE 38.6](image-url) Free precessional orbits of the free layer of a magnetic device with a $3 \times 50 \times 75 \text{nm}$, $M_s = 800 \text{kA/m}$ free layer. The free layer has a perpendicular anisotropy with an anisotropy field of $\mu_0 H_{an} = 1.25 \text{T}$. (a) With no damping and no applied field, (b) with no damping and $0.3 \text{T}$ applied field in the $x$-direction and $y$-direction, (c) with damping ($\alpha = 0.02$) and $0.3 \text{T}$ applied field in the $x$-direction and $y$-direction, and (d) with damping ($\alpha = 0.02$) and no applied field.
Spin-Transfer Nano-Oscillators

\[ f_r \equiv \frac{\gamma}{2\pi} \mu_0 H_{ap} \]

at \( \theta = 90^\circ \) and

\[ f_r \equiv \frac{\gamma}{2\pi} \mu_0 (H_{ap} + N_{az} M_f) \]

at \( \theta = 180^\circ \).

For in-plane orbits with an applied field along the \( x \)-direction the resonance frequencies are

\[ f_r \equiv \frac{\gamma}{2\pi} \mu_0 \sqrt{(H_{ap} - N_{ax} M_f + N_{az} M_f)(H_{ap} - N_{ax} M_f + N_{ay} M_f)} \]

The first term under the radical corresponds to the torques applied at point \( b \) in Figure 38.5a due to the applied field and the \( x \) and \( z \) components of the demagnetizing fields. The second term under the radical corresponds to the torques applied at point \( a \) in Figure 38.5a due to the applied field and the \( x \) and \( y \) components of the demagnetizing fields. These equations, referred to as Kittel equations (1948), illustrate a general property of magnetic devices, since the effective field depends on the magnetization vector, the frequency of the precessional orbit varies for different orbits. \(^\ddagger\) For the case shown in Figure 38.5a where there is no applied field, the frequency formula reduces to

\[ f_r \equiv \frac{\gamma}{2\pi} \mu_0 M_f \sqrt{(N_{az} - N_{ax}) (N_{ay} - N_{ax})} \]

indicating that the frequency will decrease as the orbit opens up and \( M_f \) decreases. The variation of the resonant frequency is critical to the operation of the STNO since the device will try to vary its orbit and orbit frequency to balance out the damping and spin torques that will be discussed in the next sections. Note that \( f_r \) is the fundamental frequency of the orbit and, depending on the configuration, there can be higher harmonics. For in-plane orbits, the component along the axis of the orbit \( m_{\parallel} \) will, as seen by inspecting the in-plane orbits in Figure 38.5a, oscillate at twice this frequency.

When the field is not applied along a high symmetry direction, as shown in Figure 38.5c, the orbits can have a distorted shape and there will be many higher harmonics present. This illustrates the non-linearity of the torque equation due to the dependence of the effective field on the moment direction. The magnetization components will in general not be simple functions of the form \( \cos(\omega t + \phi) \).

\(^\ddagger\) The dependence of precession frequency of a ferromagnetic element on the orbit amplitude and orientation is quite different than what is typically found in nuclear magnetic resonance (NMR) where there are no internal effective fields and the frequency of the orbit is independent of the amplitude of the orbit. This is because the torque is proportional to the transverse component of the moment times the applied field. As the orbit grows in size the torque increases with the size of the orbit thus maintaining the frequency.

An alternative geometry is to induce a perpendicular anisotropy using the correct magnetic materials. Orbits for a device with perpendicular anisotropy \( \mu_0 H_{an} = 1.25 \) T and no damping are shown in Figure 38.6a and b. When the anisotropy field is larger than the magnetization the equilibrium state of the moment points perpendicular to the plane of the device. In Figure 38.6a there is no external field applied and the moment precesses around the sum of the magnetostatic and anisotropy fields. The effective field now has the opposite sign compared to Figure 38.5a and the rotational direction is reversed. Since the anisotropy and \( z \)-component of the magnetostatic field have the same form but opposite signs, we can use the previous formulas for resonant frequencies with the substitution \((-N_{ax} M_f) \Rightarrow (H_{an} m_{\parallel} - N_{az} M_f))\.

The frequency is given by

\[ f_r \equiv \frac{\gamma}{2\pi} \mu_0 \left[ H_{ap} + (H_{an} - N_{az} M_f) \cos(\theta) \right] \]

and again, as \( m_{\parallel} = \cos(\theta) \) decreases and the orbit opens up, the orbit frequency will decrease until \( m_{\parallel} \) changes sign and then it will increase, rotating in the opposite sense. In Figure 38.6b an external field at an angle \( \theta = 45^\circ \) is applied with components \( \mu_0 H_{ap} = 0.3 \) T and \( \mu_0 H_{an} = 0.3 \) T. The larger orbits are again non-elliptical with substantial higher harmonic components.

### 38.3.2 Equation of Motion with Damping

The equation of motion discussed above assumes that the local moment does not interact with its environment except through a reversible applied field. In actuality, the moment interacts strongly with its environment, particularly the conduction electrons and phonons. This interaction leads to dissipation and the relaxation of the moment into a low energy state. To account for these interactions the equation of motion has to be modified:

\[
\frac{d\vec{m}}{dt} = -\mu_0 \frac{\gamma}{1 + \alpha} \vec{m} \times \vec{H}_{\text{eff}} - \alpha \frac{\gamma}{1 + \alpha} \vec{m} \cdot \vec{H}_{\text{eff}}
\]

This equation, referred to as the Landau Lifshitz Gilbert (LLG) equation,\(^\ddagger\) has three modifications. First, we have to add a damping term (second term on the right) that causes the moment to relax in a direction toward the effective field which moves the system toward a low energy state. The triple cross product in the damping term pulls out the component of \( \vec{m} \) to \( \vec{H}_{\text{eff}} \) that is normal to the free layer moment so that the magnitude of the moment is still conserved. The strength of the relaxation is determined by the dimensionless damping parameter \( \alpha \). Second, we have to add the factor \( 1/(1 + \alpha^2) \) to both the first and second terms. This causes viscous damping: in the limit of large \( \alpha \) the magnetization responds slowly. The third change is buried within \( \vec{H}_{\text{eff}} \).

We need to add a random thermal fluctuation field \( H_{\text{th}} \) to \( H_{\text{eff}} \).

\(^\ddagger\) For a recent comprehensive review of the LLG equation see Bertotti et al. (2008).
$H_a$ is typically chosen to be a random Gaussian-distributed field with a root-mean-square (rms) average value of

$$\mu_s H_{a,\text{rms}} = \frac{2kT\alpha}{\sqrt{VM}\gamma\Delta t},$$

where $\Delta t$ is the period over which the thermal field is applied. It is this random field that gives rise to paramagnetic response when the magnetic element becomes sufficiently small and gives the classic Langevin magnetic behavior. The damping term describes phenomenologically the average interaction of the magnetic device with its environment while the thermal field describes the stochastic part. The magnitude of the damping constant for materials used in magnetic devices is typically between 0.005 and 0.05 and is always positive.

When we incorporate damping into the equations of motion we see (Figures 38.5d and 38.6c,d) that the orbits relax slowly to the nearest energy minimum. The relaxation time is on the order of 1–5 ns. In addition, there are thermal fluctuations that cause some random motion of the orbits. Since the damping constant is typically small for materials of interest, the damping term can be considered a small perturbation to the precessional term. As the moment relaxes, it will sample a family of nearly constant energy orbits. To first approximation, these are the orbits that are stabilized by the transfer of spin from the applied current.

38.3.3 Equation of Motion with an Applied Current

John Slonczewski (1996) realized that the LLG equation would have to be modified in magnetic devices since the electron currents transport angular momentum from one layer to another

$$\frac{d\vec{m}_l}{dt} = -\mu_s \frac{\gamma}{1+\alpha^2} \vec{m}_l \times \vec{H}_{eff} - \alpha \frac{\gamma}{1+\alpha^2} \vec{m}_l \times (\vec{m}_l \times \vec{H}_{eff})$$

$$+ \frac{g(\theta_m)\mu_e I}{eM_V} \vec{m}_l \times (\vec{m}_l \times \vec{m}_l).$$

The new term accounts for the angular momentum transferred by a current $I$ from a layer with polarization direction $\vec{m}_l$, as discussed in Section 38.2. Here, we only consider the transverse torque and ignore the smaller field-like torque. The magnitude of the spin-transfer term normalized to the precessional term is approximately $g(\theta_m)\mu_e I / \mu_s e M_V^2 \approx 0.03$, where we have assumed that $H_{eff} \sim M_e$, an applied current of 1 mA, and the same device dimensions as used in Figures 38.5 and 38.6. For a device of this size the maximum current before catastrophic failures is approximately 10 mA. Since the spin-transfer term is small, on the order of the damping term, it can also be viewed as a small perturbation to the precessional term. The equation of motion with the spin-torque term is referred to as the Landau–Lifshitz–Gilbert–Slonczewski (LLGS) equation and it has been extensively studied both numerically (Sun 2000) and analytically (Bertotti et al. 2005).

Depending on the relative orientation of the polarization direction to the effective field, the spin-transfer term can either add energy or remove energy from the magnetic system. To understand how these terms affect the magnetization dynamics it is useful to look at a simple high symmetry configuration. In the case of perpendicular applied field and polarization, as shown in Figure 38.8, the LLGS equations can be rewritten as

$$\frac{d\vec{m}_l}{dt} = -\mu_s \frac{\gamma}{1+\alpha^2} \vec{m}_l \times \vec{H}_{eff} - \left( \frac{\alpha}{1+\alpha^2} \frac{g(\theta_m)\mu_e I}{eM_V} \right) \vec{m}_l \times (\vec{m}_l \times \vec{H}_{eff})$$

We can see that the spin-torque term opposes the damping term if $I$ is positive and adds to the damping if $I$ is negative. When the current becomes greater than a critical current the effective damping becomes negative and the system will move toward orbits of higher energy. The critical current is given by setting the effective damping

$$\alpha_{eff} = \frac{\alpha}{1+\alpha^2} \frac{g(\theta_m)\mu_e I}{eM_V}$$

to 0

$$I_c = \frac{\alpha}{1+\alpha^2} \frac{eM_V}{g(\theta_m)\mu_e}.$$

To obtain a small critical current to set the free layer in motion, devices need free layers with small damping constants, low saturation magnetizations, small volumes, and high polarizations. Reducing $H_{eff}$ is in general not feasible since a large $H_{eff}$ is required to maintain a high operation frequency. Conversely, when designing a fixed polarizer layer, we require a large magnetization, large damping, and large effective field so that the current is not sufficient to get the polarizer precessing. The above discussion applies strictly to the case when everything is symmetric about the $z$-axis. In general, the effective damping will not be zero everywhere on the orbit. Stable orbits are obtained by having the integrated damping and spin-torque terms balance on the orbit.

Having the effective damping go to zero is not sufficient to obtain a stable orbit. For stability we require positive damping if the orbit is perturbed to a higher energy level and negative damping if the orbit is perturbed to a lower energy level. For the case of a circular device with the applied field and polarization directions along the $z$ axis, as shown in Figure 38.8, the effective field is given by $H_{eff} = H_{ap} - N_x M_s \cos(\theta)$. This term gives rise to nonlinear damping and causes the damping to increase when the orbits increase in size. When the current is increased to make the effective damping go negative the orbit begins to grow, $\theta$ increases, and the effective damping increases until it becomes zero. If the orbit grows beyond this point the damping goes positive pushing the orbit back to the $0$-damping orbit. In this geometry, the LLGS equation becomes similar to the
van der Pol equation that has non-linear damping. A STNO, similar to a van der Pol oscillator, is an auto-oscillator that has a resonator driven by a nonlinear damping term that adjusts the amount of energy being put in and taken out of the system to stabilize a particular oscillation mode. It should be pointed out that, unlike the van der Pol equation, the LLGS equation has a non-linear precessional term so that the oscillation frequency, as seen in Figures 38.5 and 38.6, varies depending on oscillation amplitude.

It is instructive to write down the $z$-component of the LLGS equation in this geometry

$$
\frac{dm_z}{dt} = \left( \alpha \|H_{ap} - N_{m} M_s \cos(\theta) \| + \frac{g_s(\theta_m)\mu_I I}{cM_s V} \right) \sin^2(\theta).
$$

Since the terms in this equation are small the $z$-component of the magnetization, unlike the oscillating transverse components, will evolve slowly. The right side of the equation is plotted as a function of the polar angle in Figure 38.7 for a $z$-axis applied field of 1.25 T and typical STNO parameter values. If $\frac{dm_z}{dt} > 0$ then the orbit is pushed toward smaller angles, if $\frac{dm_z}{dt} < 0$ the orbit is pushed toward larger angles and if $\frac{dm_z}{dt} = 0$ than stationary solutions are found. If the current is zero, the moment is pushed to a static state at $\theta = 0$ as expected. As the current is increased, this point becomes unstable and the system will evolve to a stable precessional state. At large currents the system will be pushed to a static state at $\theta = 180^\circ$, i.e., the device switches and no stable dynamics are observed. The energy of the STNO is also plotted in Figure 38.7 as a function of the free-layer polar angle. Since the energy can wander by several $k_BT$ due to thermal fluctuations the orbit angle will also wander giving rise to variations of the STNO frequency and a finite linewidth. If the precessional state has a large slope at the zero crossing then the restoring force for the orbit will be large and a sharp emission peak will be observed. If the slope at the zero crossing is small, then the orbit will have a small restoring force and the frequency will tend to wander.

The spin-torque parameter, $g_s$, is also a function of angle if $\Lambda \neq 1$, increasing monotonically as the angle $\theta_m$ between the free layer and polarization layer moments varies from 0° to 180°. This dependence will make the system asymmetric: inverting both the sign of the current and orientation of the polarizer will not give the same orbits. Starting with the free layer moment parallel to the polarizer will cause $g_s(\theta_m)$ to increase as the orbit opens up and may prohibit stable orbits, while starting with the free layer magnetization antiparallel to the polarizer will cause a decreasing $g_s(\theta_m)$ which will stabilize the orbits. Other device parameters, such as the damping coefficient $\alpha$, may also be a function of the magnetization angle. Further, the spin transfer and damping torques will, in general, be along different directions with varying amplitudes making a general analysis of orbit stability difficult.

Examples of current stabilized orbits are shown in Figure 38.8a for a device with a perpendicular applied field, perpendicular polarizer, and $\Lambda = 2$. The orbits shown are numerically calculated from the LLGS equation with an applied field of $-1.5$ T so that the magnetization starts antiparallel to the polarizer and the current is positive, which pushes the free layer moment toward the polarizer. This orientation has the largest region of stable precession. Figure 38.8c, d shows “phase diagrams” of the steady state response of the system to different applied fields and currents. Figure 38.8c plots the average $rms$ value of $m_z$ using a color scale: blue is zero and indicates a quiescent state, red is 0.7 represent a state of maximum oscillation. Figure 38.8d plots the average value of $m_z$ which show the orientation of the magnetization as a function of applied field and current. For a given field value, as the current is ramped, the magnetization goes from being aligned with the polarizer’s direction at negative currents to being anti-aligned with the polarizer at large positive currents. This magnetization rotation is reflected in a change in device resistance from its minimum value to its maximum value. A region of oscillation is often identified by a region of intermediate device resistance.

The high-symmetry configurations, while being useful to understand how STNOs function, often do not make practical STNOs. In the case shown in Figure 38.8, the angle $\theta_m$ between the free layer and polarizer remains constant on the precessional orbit so there is no time-varying voltage output. A third magnetic reference layer is required to cause a changing resistance and an output voltage.

It can be seen in Figure 38.8 that the orbits are not perfectly sharp and, in fact, there are dynamical excitations present when the current is zero. This is due to the thermal field which continually agitates the system and gives rise to a finite linewidth as discussed in Section 38.7.
Many different STNO configurations have been explored. The configurations, as seen in Figure 38.9, can be classified by the type of spacer layer, the patterning geometry, and the magnetic geometry. The main requirements are that there is efficient generation of spin torque to induce oscillations without damaging the device with large electrical currents and there is a large magnetoresistance in the correct geometry to convert the time dependent magnetization into a useful output voltage.

### 38.4.1 Barrier Type and Stack Configuration

Devices that have metallic spacer layers, as shown in Figure 38.9a, are referred to as spin valves and those with thin insulating layers, as shown in Figure 38.9b, are called MTJs. The term "spin valve" was coined by Dieny et al. (1991) and refers to a system with two magnetic layers separated by a metallic spacer that lets spins through when the moments are aligned and blocks electrons when the layer magnetizations are anti-aligned.

A typical stack is shown in Figure 38.9a and consists of a 2.5 nm Ta adhesion layer, a 50 nm Cu layer that serves as a high conductivity base electrode and a seed layer that nucleates (111) texture, a 20 nm Co_{90}Fe_{10} layer that serves as the spin polarizer, a 5 nm Cu spacer layer, a low moment 5 nm Ni_{80}Fe_{20} free layer, and a top Cu capping layer. The free layer is made from a magnetically soft Permalloy alloy that can easily respond to both applied fields and spin torques. The polarizing layer, in this design, is thicker and has a higher moment than the free layer to prevent it from undergoing spin-torque induced oscillations. CoFe alloys are used to obtain fcc texture which is consistent with the desired (111) texture in the stack. Pure Co is likely to grow in an hexagonal close packed (hcp) structure which is undesirable for these devices.

All metallic devices, such as spin valves, have low resistances on the order of 2–5 \( \Omega \). Accurate measurement of the intrinsic impedance of an all metallic device is difficult since the resistance of the electrodes and contacts contribute a considerable fraction of the device resistance. A typical device may have a \( \Delta R = 0.2 \Omega \), intrinsic resistance of \( R_{int} \approx 1 \Omega \), and a total resistance of \( R_{tot} = 5–10 \Omega \). The devices have a low resistance area product of \( RA \approx 10^{-2} \Omega \mu m^2 \). Given the low resistance of all
metallic devices, it is easy to apply large currents to initiate spin-transfer oscillations. However, since the resistance change is relatively small the output power is also relatively small. The maximum observed output powers range from 1 to 10 nW.

Tunneling devices have thin insulating barriers. The most successful barrier material to date is MgO (Yuasa et al. 2004). For STNO applications, the junctions must have a high MR and a high transparency (low RA product). The high transparency is required because the TMR and spin torque fall off at voltages above a few hundred mV and the junction breakdown voltages are typically 0.5–1 V. To safely get 1 mA, a typical critical current, through a 100 nm junction requires RA < (1 V · 10^{-2} μm^2/1 mA) = 10 Ω μm^2. This RA product corresponds to MgO barriers that are less than 0.7 nm thick. To obtain a high polarization and ΔR, the MgO barrier requires a (001) crystalline texture which is achieved by annealing the MTJ stack at high temperatures (300°C–400°C). MTJs have the advantage that ΔR and R_{av} can be much larger, typically 100 and 200 Ω, respectively.

In addition to metallic and MTJ devices, a third class of devices has been developed using porous nano-oxide films as the spacer layer. These devices can be viewed as an intermediate structure between a MTJ and a metallic device and has the advantage of a higher device resistance than an all metallic device and it can accommodate more current than an MTJ.

The stack structure for either a spin valve or a MTJ can be a simple bilayer, as shown in Figure 38.9a or a pinned bilayer as shown in Figure 38.9b. The simple bilayer usually has a fixed layer that is thicker or has a higher magnetization material to stabilize the layer to prevent it from oscillating. A pinned stack uses an antiferromagnetic layer (usually Ir_{50}Mn_{50} or Pt_{50}Mn_{50}) and a synthetic antiferromagnetic (SAF) to fix the direction of the polarizer. The antiferromagnetic layer couples to the adjacent ferromagnetic layer (pinned layer) through the exchange interaction. Since the antiferromagnet has no moment it does not respond to an applied field and it is difficult to rotate the magnetic order. The direction of the pinned ferromagnetic layer is set by depositing or annealing in a magnetic field. This orients the pinned layer in the desired direction and then allows the antiferromagnet to relax into a low energy state in which the antiferromagnetic moments line up with the ferromagnetic moments. When the ferromagnetic moment attempts to rotate away from the preferred direction, it experiences an exchange bias field that tries to keep it in its pinned direction. The exchange bias is an interfacial energy proportional to the area of contact whereas the Zeeman energy is proportional to the volume or total moment of the magnetic element. This causes the exchange bias field to vary inversely with the moment of the layer so thin layers are required to get large pinning fields.

A SAF is often used for a pinned layer instead of a single ferromagnetic layer. A SAF consists of two ferromagnetic layers separated by a thin layer of Ru, which promotes strong antiparallel orientation. The effective moment of the SAF is the difference between the two layer moments. For the device shown in Figure 38.9b, the effective moment of the SAF is close to zero and the pinning strength can be very high 0.1–0.5 T.

The device stacks are usually deposited in one step using a computerized sputter system with 6–10 cathodes. The sputter system should have base pressures of <10^{-6} Pa (0.75 × 10^{-8} torr)

FIGURE 38.9 Different configurations of STNOs based on the type of barrier, the patterned geometry, and the magnetic state.
since small amounts of residual gases can affect the microstructure of the various layers. The deposition conditions, such as deposition pressure and power, need to be carefully controlled to ensure that the interfaces are sharp and flat, the crystalline texture is correct, and the magnetic properties are optimized. Often, small partial pressures of reactive gases (N₂ or O₂) may be used during one or more deposition steps to control the film microstructure.

**38.4.2 Patterning Geometry and Magnetic Configuration**

There are three basic patterning geometries: nanopillars in which the magnetic structure is patterned down to nm dimensions (Figure 38.9c), half patterned nanopillars in which the polarizer is left unpatterned (Figure 38.9d), and nanocontacts in which a small <100 nm contact is made to a larger mesa (Figure 38.9e). The magnetic structure of a nanopillar is very sensitive to the shape of the element and there is a strong interaction between the free layer and polarizer. A half patterned nanopillar reduces the interaction between the two layers and makes the polarizer more likely to be fixed in a magnetic field since it has a higher volume. The nanocontacts have the advantage that there are no device edges to induce large magnetostatic fields that may inhibit the ability of the device magnetization to precess uniformly. However, as will be seen in the next section, the exchange coupling of the oscillating region to the larger mesa gives rise to spinwaves.

Each of these geometries can use a variety of magnetic orientations with the free and fixed layer being designed to be in-plane or perpendicular to the plane. Figure 38.9f shows a simple in-plane device where the magnetizations like to lie in plane due to the magnetostatic fields. Figure 38.9g shows a device with a perpendicular free layer, which is shown schematically as a multilayer. The fixed layer still lies in plane. This is an efficient geometry since it is easy to get the free layer to precess and there is a large angular variation between the free layer and the polarizer leading to a large MR and output voltage. Figure 38.9h shows a geometry in which both the free layer and polarizer are perpendicular. More advanced devices can have three or more magnetic layers and can use complex bit shaping to obtain different modes of operation.

Etching of STNO devices still remains challenging due to the large number of different materials in the devices stacks and the lack of selective etches. Etching is usually done with an ion beam etch or a reactive ion etch process. Care must be taken to prevent redeposition on the devices’ edges during etching or oxidation of the device edges as this can alter both the electron conduction though the device and the magnetic properties of the free layer. The nanocontact geometry avoids many of the complexities involved with the etch process required for nanopillar geometry.

**38.5 Dependence of Frequency on Current and Field**

STNO frequencies vary with both applied field and current. As the magnetic field is increased, the precession frequency will generally increase as the square root of the applied field at low fields and linearly at high fields as predicted by the Kittel equation. An example of the frequency and field variation of a spin valve nanocontact oscillator with in-plane applied fields is shown in Figure 38.10, along with the predicted single domain orbits from the single domain model with 60 mT applied field.

---

**FIGURE 38.10** Measured frequency and field dependence of a spin valve nanocontact oscillator (Rippard et al. 2004) with in plane applied fields (a) field dependence, (b) output spectra for various applied currents with 60 mT applied field showing first and second harmonics, (c) predicted orbits from the single domain model with 60 mT applied field.
orbits. A red shift (frequency decrease) is seen as the current is increased as expected from the discussion in Section 38.3. A strong second harmonic signal is observed since the output voltage is predominantly proportional to $m_f x$.

As the applied field is moved out of plane, we expect to see spectra that red shift then blue shift (increase frequency) with current, as seen in Figure 38.5c, and then at large, nearly out-of-plane fields we expect to see solely a blue shift, as seen in Figure 38.5b. This trend can be seen in the data shown in Figure 38.11, which plots the oscillator frequency as a function of applied current and field for a spin valve nanocontact oscillator when the magnetic field is applied 10° off normal. At low fields, one sees a slight red shift then blue shift as the current is increased. At high fields, there is a strong blue shift with an average magnitude of $\pm 1.5 \text{GHz/mA}$. There are, however, sharp jumps in the frequency suggesting that the device abruptly changes oscillation modes at certain fields and currents. These modal jumps have not yet been fully explained. Figure 38.11 shows data for both increasing and decreasing currents highlighting the fact that, for magnetic devices, there can be hysteresis and the behavior of a device may depend on its current and field history. Also shown in the color maps are the STNO linewidth (Figure 38.11a) and power output (Figure 38.11b). For spin-valve nanocontact STNOs the linewidths vary from a few MHz to several hundred MHz and the maximum output power is on the order of 1 nW. Output power and linewidths will be discussed in more detail in the next sections.

The frequency output from a spin valve nanocontact STNO with a perpendicular free layer and an in-plane polarizer is shown in Figure 38.12a. Here $\mu_s(H_{an} - M_s) \equiv 0.20 \text{T}$. As previously discussed, for this geometry, we expect a red shift as the current increases (as seen in Figure 38.6a). The moment initially points vertically at low currents with no precession. As the current is increased, the orbit should open up and the effective field $H_{ap} + (H_{an} - M_s) \cos(\theta)$ will decrease causing the frequency to decrease. The precessional angle, shown in Figure 38.12b, can be calculated from the data knowing the saturation magnetization and the perpendicular anisotropy field. In this configuration with an in-plane polarizer, the orbit at larger currents will saturate near 90° and will begin to deviate from the “free-precession” orbits. The maximum output power of 1 nW is, as expected, observed near 90° when the orbit obtains its maximum amplitude. One advantage of this type of perpendicular oscillator is that it can operate with little or no applied field.
38.6 Output, Circuit Model, and Device Measurement

For an oscillator to be useful there needs to be an output. There are two possible outputs generated by an STNO: (1) a microwave voltage and (2) strong oscillations in the magnetic field close to the oscillator. Using the strong local microwave magnetic fields for data recording applications has been proposed by Zhu et al. (2008). The microwave fields within a few tens of nanometers from the device can be on the order of a 1–2 T. There are very few devices that are capable of producing such large microwave magnetic fields. These fields can resonantly drive magnetic data storage bits to assist in their switching.

More conventionally, we are interested in voltage and current outputs. A STNO may be viewed as a time dependent resistor as shown in Figure 38.13. The STNO resistance is given by

\[ R = R_{av} + \frac{\Delta R}{2}(\overline{m}_I \cdot \overline{m}_p) \leq R_{av} + \frac{\Delta R}{2}\cos(\omega t). \]

In Figure 38.13, we have the STNO connected through a bias tee to a load \( R_L \). We assume the inductor is sufficiently large and close to the device so that the drive current remains constant. The changing STNO resistance alternately directs the current through the STNO when its resistance is low or through the load when its resistance is high. The optimal device would be one that would change its resistance from zero to an infinitely large value when its resistance is high. The optimal device would be one that can get 16 \( \mu \)W. Of course, in both cases we have increased the drive power by a factor of 10 and have increased the output power by slightly more than this because we have impedance matched to the load. The efficiency

\[ \frac{P_{out}}{P_{in}} = \frac{V_{out}^2}{2IR_L} = \frac{1}{8} \frac{\Delta R^2 R_L}{(R_{av} + R_L)^2} \]

in both cases is low, ranging from \( 0.5 \times 10^{-3} \) to \( 5 \times 10^{-5} \). This can be compared to the efficiency of Gunn oscillators which is typically 0.03–0.05. Strategies to increase the efficiency will be discussed in Section 38.9, which discusses phase locking.

All devices will have a parallel shunt capacitance, as seen in Figure 38.13, that could potentially limit high-frequency operation. Here again, nanoscale dimensions come to the rescue and intrinsic parallel capacitances are on the order of 1 aF, which will allow operation up to 300 GHz. Care still must be taken when designing the circuit layout to avoid parasitic capacitances from the electrodes.

38.7 Linewidth

Two of the most important parameters characterizing an STNO are the oscillator’s full width at half maximum (FWHM) linewidth \( \Delta f \) of the power spectra and the quality factor \( Q = f_r / \Delta f \). Figure 38.14 shows an STNO emission spectra with \( f_r = 13.15 \text{ GHz}, \Delta f = 7.1 \text{ MHz}, \) and \( Q = 1800 \). The linewidth can vary widely depending on the device and operation conditions. As shown in Figure 38.11a using the color scale, the linewidth for a single device can vary from a few MHz to 100 MHz as the current and applied field are varied and the linewidth variation can be very complex. One trend that can be observed in Figure
The Fourier transforms of $m_z$ is plotted and the linewidths listed are the half width at half maximum which correspond to the power spectra FWHM linewidths. (b) Calculated temperature dependence of the STNO linewidth along with a fit which shows a linear behavior at high temperatures. The device configuration is shown in the inset.

38.11a is that the linewidth tends to be larger near regions of rapid frequency variation and smaller in plateau regions where the variation of frequency is small. In general, linewidths for nanocontacts are smaller than that for nanopillars, where linewidths can often approach 1 GHz.

One contribution to the observed linewidth is due to thermal fluctuations that perturb the amplitude and phase of the orbit. LLGS simulations of STNO emission at $T = 30$, 100, and 300 K are shown in Figure 38.15 along with an inset showing the thermally perturbed trajectories. The Fourier transform of $m_z$ is plotted, which is proportional to the output voltage since, for this simulation, the polarization is predominantly along the $x$ direction. The peak width for the voltage spectra is twice that for the power spectra and therefore half width at half maximum are listed to correspond to the power spectra widths. The calculated linewidth is approximately linear in temperature with the linewidth going to zero at $T = 0$ K.

Experimentally, STNO devices show a decrease in linewidth with temperature, however, the linewidth tends to saturate at temperatures near 100 K. The linewidth may be limited by micromagnetic effects caused by devices edges, defects, or inhomogeneous current injection that cause different parts of the device to oscillate at slightly different frequencies. Experimentally, many parameters vary with temperature, including $M_s$ and $\alpha$, causing the orbit to change as a function of temperature. As indicated in Section 38.3, the linewidth is sensitive to the stability of the particular orbit and different orbits can have widely varying linewidths. This makes the analysis of the temperature dependence of the linewidth difficult. As indicated, point contact devices typically show smaller linewidths than nanopillar devices. This may be due to a larger effective area (the coherent oscillating region is larger than that of the contact) and reduced inhomogeneities near the region of oscillation that reduce the inhomogeneous linewidth broadening.

Most of the oscillator linewidth is due to phase noise, the random deviation of the oscillator phase from that of an ideal reference oscillator. The amplitude noise is relatively small, however, since the amplitude of oscillation changes the frequency, the amplitude and phase noise are coupled. In addition to thermal magnetic fluctuations, which are associated with the damping constant $\alpha$, there are also Johnson current fluctuations associated with the device resistance $R$. The Johnson current noise is typically small in metallic devices while, in MTJs both Johnson noise and shot noise may contribute to the STNO linewidth. These noise sources explain the correlation between the observed linewidth and the change in frequency with current and field: if the frequency changes rapidly with field and current then fluctuations in these quantities will cause considerable phase noise and an increased linewidth.

38.8 Micromagnetics, Vortex Oscillators, and Spinwaves

In the previous sections, we considered the magnetic system to consist of a uniformly magnetized free layer and a fixed polarization layer. These assumptions are not always justified. The magnetization in both (or several) layers may be nonuniform and both layers can undergo dynamic response. The nonuniformity can be driven by the magnetostatic energy which causes the magnetization to rotate at the device edges and the current induced magnetic fields (sometimes called the Oersted fields) that promotes a vortex structure. The LLGS equations of motion can be generalized to accommodate spatially nonuniform magnetization and dynamics in multiple layers by discretizing the system, as shown in Figure 38.16, and solving a large set of coupled integro/differential equations.
An example of micromagnetic effects is shown in Figure 38.17 where two nanocontact STNOs are oscillating on the same magnetic mesa. When the region under the contacts oscillates, spin waves are generated that cause the oscillators to interact. The spin waves can be detected directly by using one STNO as an emitter and one as a detector (Pufall et al. 2006). When both STNOs are oscillating and the frequencies become close the two oscillators can phase lock. Figure 38.17 shows the mesa structure after a line was cut by a focused ion beam which then causes the spinwave coupling to disappear. The simulation shown in Figure 38.17 does not account for the Oersted fields generating by the currents. These fields can cause the radiation pattern to be very asymmetric. The detailed spin wave radiation pattern generated from these nanocontact devices has not yet been experimentally determined.

Another example of micromagnetic effects is the vortex oscillator. A vortex oscillator in the nanocontact geometry, after Pufall et al. (2007), is shown in Figure 38.18. By adjusting the thickness of the magnetic layers and going to large device currents, which create large circumferential magnetic fields, vortices can be formed in either the free layer or the polarizer layer (Pribiag 2007). A vortex has a core in which the moment points out of the plane of the device and the magnetization surrounding the core is in a circumferential in-plane direction. The spin torque causes the vortex to move around the region of current injection. Vortex oscillators usually have low frequencies near...
1 GHz and show little dependence of the frequency on drive current. The oscillator shown in Figure 38.18 is highly nonlinear with large harmonic amplitudes indicating that the vortex is not undergoing simple harmonic motion. One advantage of the vortex oscillators is that they require small or no applied magnetic field for operation.

There are many further examples of using nonuniform dynamical states in STNOs. The reader is referred to Silva and Rippard (2008) for more complete discussion of micromagnetic effects.

### 38.9 Phase Locking

An important aspect of an oscillator is to control the phase of the oscillator with respect to a reference oscillator or other oscillators in an array. This is required if one wants to adjust the phase of an oscillator in a phased array antenna or to lock a set of oscillators together to enhance the total power output. It is a general property of non-linear auto-oscillators that the oscillator will try to adjust its frequency to match that of an injected signal. In the case of STNOs the injected signal can be either a microwave current or magnetic field. Figure 38.19 shows an example of phase locking to an injected microwave current. The device geometry is shown in the inset in Figure 38.19a. With no applied current the frequency increase with applied current as shown in Figure 38.19a. When a microwave current is injected, the oscillator will lock over a given range of applied current. This locking range is inversely dependent on the linewidth of the oscillator and is typically 0.5–2 mA. When the signal is locked, there will be a low frequency voltage that appears, seen in Figure 38.19f, due to the mixing of the applied microwave signal with the oscillator output. When the oscillators are locked, the frequencies will be the same but there will be a phase difference between the two oscillators that varies over the locking range. This property is useful to adjust the phase of the STNO relative to a reference oscillator by adjusting the bias current.

Another important example of phase locking is the coupling of two or more STNOs to have them oscillate at the same frequency with an adjustable relative phase. An example of mutual phase locking of two point contact STNOs is shown in Figure 38.20 (Kaka et al. 2005, Mancoff 2005). The geometry is similar to that shown in Figure 38.15. When the oscillators are far apart, the frequency of the oscillators can be swept and there is little interaction between them. In the experiment shown, the current through oscillator B was swept and oscillator A was kept at a fixed current of 11.5 mA. As the oscillators are brought close in space, 500 nm as shown in Figure 38.20b, and frequency they interact and phase lock. In this case the phase locking is due to a spin-wave coupling between the oscillators. Figure 38.21 shows the spectral outputs before, during, and after phase locking. When the oscillators are locked, the frequencies will be the same but there will be a phase difference between the two oscillators that varies over the locking range. This property is useful to adjust the phase of the STNO relative to a reference oscillator by adjusting the bias current.

![Image of phase locking](image-url)

**FIGURE 38.19** (See color insert following page xxx.) Measurement of an STNO phase locking to an injected microwave signal. The device geometry is shown in the inset (a); (a), (b) and (c) show the measured frequency, device resistance, and output power as a function of bias current with no injected microwave signal. (d), (e), and (f) show the measured frequency, device resistance, and dc voltage as a function of bias current with an injected microwave signal. (Adapted from Rippard, W.H. et al., Phys. Rev. Lett., 95, 067203, 2005.)
FIGURE 38.20  (See color insert following page xxx.) Demonstration of mutual phase locking of two STNOs (Kaka et al. 2005). (a) The emission spectra of two oscillators as a function of the current through oscillator B when the oscillators are far apart showing no interaction. (b) The emission spectra of two oscillators as a function of the current through oscillator B when the oscillators are close showing strong interactions and phase locking.

FIGURE 38.21  Spectral output of a two element STNO array (a) when STNO B has a frequency less than STNO A, (b) when the oscillators are locked showing an increase in output power and decrease in linewidth and (c) when STNO B has a frequency greater than STNO A, (d) the output power as a function of a relative phase shift before the signals are combined. (From Kaka, S. et al., Nature, 437, 389, 2005.)
38.10 Applications

Several applications have been suggested for STNOs. Since power output, in general, scales with the size of the oscillator and the linewidth inversely proportional to the size, a single nano-oscillator will never be a high power ultra-narrow-linewidth oscillator. Making phase locked arrays may solve this problem if high power is required. A better set of applications are those that make use the nanoscale aspects and large-range high-speed tunability of STNOs. These applications include local on chip clocks for VLSI applications, high-density massively-parallel microwave signal processors, small phased array transmitters, chip-to-chip micro-wireless communications, and local excitation sources for nanosensors. Figure 38.22 illustrates an array of STNOs being used for rapid demodulation of an incoming signal. This application relies on a large array of variable frequency STNO demodulators working simultaneously to analyze an incoming signal. Figure 38.23 illustrates a nanoscale STNO phased array antenna that could have possible applications in chip-to-chip communications. The relative phase of each oscillator can be adjusted

![Figure 38.22](image1)

**FIGURE 38.22** Schematic showing a possible application using an array of STNOs for rapid signal processing of an incoming signal. The application is enabled by the ability to have a large number of STNOs on the same chip each operating at a different frequency.

![Figure 38.23](image2)

**FIGURE 38.23** Schematic of a phased array antenna using STNOs. The ability to have a large number of oscillators with precisely controlled phase may allow beam steering in micro-wireless applications. The inset on the lower left shows the phase of an STNO as the current is varied through its locking range.
by varying the bias current as seen in the plot in the lower left. The use of a small phased array fabricated on a standard CMOS wafer would allow efficient beaming of transmit and receive signals for a local microwave communication bus.

38.11 Summary and Outlook

STNOs are exciting new nanoscale spin-based devices that offer many unique capabilities: nanoscale size, wide frequency range, high tunability, operation over a wide temperature range, nonlinear operation, and compatibility with standard CMOS processing. Many challenges remain before these devices can be incorporated in nanotechnologies. Experimentally the demonstrated frequency range is 0.5–250 GHz;* however, to obtain high frequencies, large magnetic fields have been required, which is not practical for applications. The highest frequency that has been demonstrated without an applied field is ~5 GHz. Several strategies exist to boost the frequency in the absence of an applied field. These include introducing stronger magnetic anisotropies, exchange coupling to a fixed layer, and using antiferromagnets as the oscillation layer. At present the maximum power output from STNOs is on the order of a microwatt. The power needs to be boosted into the mW range for many applications, although, for local clocks and nanoscale signal processing the existing power levels may be sufficient. There are still problems with device-to-device reproducibility, controlling the oscillator linewidths, and understanding the detailed mode structure of the oscillator. It is anticipated that most of these problems will be solved in the next few years and that STNOs will become a useful nanoscale device for a variety of applications.

References

Houssameddine, D. 2008. Spin transfer induced coherent microwave emission with large power from nanoscale MgO tunnel junctions, Appl. Phys. Lett. 93, 022505.
Rippard, W. H. et al. 2004. Direct-current induced dynamics in Co$_{80}$Fe$_{10}$/Ni$_{80}$Fe$_{20}$ point contacts, Phys. Rev. Lett. 92, 027201.

* High-frequency operation (100–250 GHz) of STNOs can be inferred from features in the high-field differential resistance data (Tsoi et al. 1998) although direct detection of microwave emission above 60 GHz has not yet been demonstrated.


**Author Queries**

[AQ1] Please check the source line of Figures 38.11, 38.19, and 38.21.


[AQ3] Please provide in-text citation for the following References: Bertotti et al. (2008), Grünberg (2008), Houssameddine (2008), Kittel (1948), O’Handley (2000).