# Perpendicular spin torque in circularly exchange-biased trilayer structures

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It is shown that both sign and magnitude of the perpendicular spin torque in magnetic tunnel junctions can be determined as a function of bias voltage by measuring the lowest eigenfrequency of a circularly exchange-biased system. A simple model allows for a qualitative and quantitative analysis.

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## I. INTRODUCTION

The prediction by Slonczewski[1] and Berger[2] that a spin-polarized current can exert a magnetic torque on layers with noncollinear magnetization has led to intense research over the past decade. This is partly due to the potential application in magnetic random access memory structures,[3] where the spin torque effect can be used to switch the magnetization direction of the storage layer relative to a reference layer, thus switching the resistance between a high-resistance state with antiparallel magnetization directions and a low-resistance state with parallel magnetization directions. Zhang *et al.*[4] pointed out that the spin torque has not only the in-plane component predicted by Slonczewski[1] and Berger[2] but also a perpendicular component. This latter component is small in metallic systems,[5–7] such as CoFe/Cu/CoFe trilayers, but significant in magnetic tunnel junctions (MTJs). Various works have presented calculations of the magnitude and bias-voltage dependence of the perpendicular component.[8–10] Experimentally, Sankey *et al.*[11] measured the line shape of the resonance induced by an rf current supplied on top of a dc bias current as a function of dc bias. By comparing their results with a macrospin model, they concluded that the perpendicular torque is quadratic in bias voltage, while the in-plane torque is linear in biasing voltage, for small voltages with magnitude less than about 400 mV. Kubota *et al.*,[12] using similar techniques, also concluded that the perpendicular spin torque is quadratic in biasing voltage. Petit *et al.*[13] measured the shift in resonance frequency as a function of bias current in CoFe/$Al_2O_3$/CoFe MTJs and concluded from the observed shift that the effective field due to the perpendicular torque is proportional to the biasing current density and thus changes sign with bias voltage. More recently, Wang *et al.*[14] refined the technique used in Ref. 11 and showed that the discrepancy between the results obtained by Sankey *et al.*[11] and by Kubota *et al.*[12] largely disappears when a more extensive data analysis is performed, which includes more contributions to the line shape and the perpendicular torque is quadratic in applied voltage. The results so obtained agree reasonably well with first-principle calculations by Heiliger and Stiles.[15] In addition, Oh *et al.*[16] have showed that in asymmetric junctions, there appears a linear dependence of the perpendicular torque on the applied voltage.

On the other hand, Li *et al.*[17] concluded that the perpendicular torque is proportional to the product of current density and biasing voltage, and that the sign of the perpendicular torque changes sign when the biasing voltage changes sign, in qualitative agreement with Petit *et al.*[13] Li *et al.*[17] extracted the perpendicular torque by measuring the switching current of MgO MTJs and by carefully accounting for heating effects due to the rather large current densities required for switching. However, since the measurements by Li *et al.*[17] are based on switching currents, the voltage range was necessarily higher than those of Sankey *et al.*[11] and of Kubota *et al.*[12] approximately 0.5–1.0 V.

Here, I suggest a different way to determine the perpendicular spin torque both qualitatively *and* quantitatively as a function of bias voltage in a trilayer system. This approach is complementary to those explored by Sankey *et al.*,[11] Kubota *et al.*,[12] and by Li *et al.*,[17] and may provide some additional insight into the perpendicular spin torque in MTJs. The idea explored here is to study the lowest-lying eigenmode of a particular MTJ trilayer system. As a bias voltage is increased, but kept small enough that the system does not become unstable, the frequency of the lowest-lying eigenmode changes linearly with the effective field due to perpendicular spin torque while the line width changes linearly with the effective field due to the in-plane spin torque. By studying the energy and linewidth dispersion as functions of applied bias voltage, the effective fields due to the perpendicular and in-plane spin torques can be determined, and their bias-voltage dependences extracted. A simple, but accurate, model allows for a quantitative analysis.

## II. EIGENMODE MODEL

The system under consideration here is a magnetic tunnel junction, consisting of two ferromagnetic (FM) layers separated by an insulating spacer, such as MgO or $Al_2O_3$, with one of the FM layers, the pinned layer (PL) exchange biased by coupling to an antiferromagnetic (AF) layer. The FM layers are taken to be appropriate soft alloys consistent with high tunneling magnetoresistance in MgO or $Al_2O_3$ MTJs. Each FM layer is disk shaped with a radius $R$ and thickness $d$, chosen such that the lowest-energy magnetic configuration within each layer is a vortex. This is the case for $R$ in the range of a few hundred nanometers to a few micrometers and $d$ in the range of a few to some tens of nanometer. The ferromagnetic exchange coupling between the two FM layers is set to zero and instead, the two FM layers interact via an interlayer exchange coupling field whose strength can be varied to simulate, for example, the spacer thickness or composition. An AF pinning layer is represented by an exchange

 

bias field, the strength of which is chosen to be consistent with typical AF/FM exchange bias systems, e.g., IrMn/CoFe, and which exerts a circular exchange[18] bias on one of the FM layers, which will be referred to as the PL. The second FM layer experiences no exchange bias from the AF layer and will be referred to as the free layer (FL). The equilibrium magnetic configuration of the two layers is then one in which both layers are in a vortex state with a common vortex chirality. It is also assumed that electrical contacts are made to the top FL and to the AF layer and that a constant biasing voltage $V_b$ applied, small enough that the response of the system is in the linear regime. A simple analytical model[19] is used to study the coupled dynamic magnetic response of the lowest-order eigenmodes excited by a magnetic field applied normal to the layers and in the presence of a bias voltage. This model describes this system as two circular disks with magnetization density $M_S$. The separation between the disks is assumed to be negligible and the magnetization in each disk is uniform along the $z$ axis, which is perpendicular to the plane of the disks. The PL is subjected to an exchange bias field $H_{eb}=H_{eb}\hat{\varphi}$ and the chirality of the magnetization of both layers is the same. The two layers interact through an effective interlayer exchange coupling field, $H_{IEC}$, arising because of interlayer exchange coupling, such as Néel orange peel coupling or interlayer Ruderman-Kittel-Kasuya-Yosida interactions. Linear perturbations are considered about the equilibrium configurations with the assumptions that the vortex cores can be neglected and that the magnetization excitations in each layer are dominated by the (1, 0) eigenmodes, which are circularly symmetric with one node at the edge of the disk in the radial direction. The radial part of these eigenmodes is taken to be $J_1(k_R r)$, with $k_R R$ the location of the first zero of $J_1(k_R R)$. These trial functions are taken to be approximate eigenfunctions[20] of the radial component of the magnetostatic field operator $\hat{h}_r[\boldsymbol{M}(\boldsymbol{r})]$ and satisfy the boundary condition of vanishing radial eigenmode on the edge of the disk. Finally, it is assumed that the magnetostatic field in the $z$ direction due to perturbations of the magnetization in layer $i=1$ and 2, corresponding to the PL and the FL, does not extend outside layer $i$ (this is based on the assumption that the magnetization density was uniform along the $z$ axis and that the thickness of the layers is very small compared to the radius of the disk).

The magnetization dynamics are described by the coupled Landau-Lifshitz-Gilbert (LLG) equations,

$$\frac{d\boldsymbol{m}_i}{dt} = -\frac{|\gamma_e|}{1+\alpha^2}[\boldsymbol{m}_i \times \boldsymbol{H}_i] - \frac{|\gamma_e|\alpha}{1+\alpha^2}\boldsymbol{m}_i \times [\boldsymbol{m}_i \times \boldsymbol{H}_i], \quad (1)$$

where $\boldsymbol{m}_i=\boldsymbol{M}_i/M_S$, $i=1$ and 2, $\alpha$ is the dimensionless damping constant, and $\boldsymbol{H}_i$ are the effective fields acting on layer $i$. The gyromagnetic factor is $\gamma_e$ and the cgs system of units is used. Equation (1) expresses the fact that the effective field exerts torques on the magnetization density. The first term on the right-hand side of Eq. (1) is the dissipationless precessional torque while the second term is the damping torque that strives to align the magnetization density with the local effective field.

By linearizing the magnetization directors about their equilibrium configurations $\boldsymbol{m}_1=\boldsymbol{m}_2=\hat{\varphi}$ and using the effective fields

$$\boldsymbol{H}_1 = H_{IEC}\boldsymbol{m}_2 - H_d m_{1z}\hat{z} - H_{eb}\hat{\varphi} - M_S N_r(m_{1r}+m_{2r})\hat{r},$$

$$\boldsymbol{H}_2 = H_{IEC}\boldsymbol{m}_1 - H_d m_{2z}\hat{z} - M_S N_r(m_{1r}+m_{2r})\hat{r}, \quad (2)$$

with $\hat{h}_r[M_S\boldsymbol{m}_{ir}(\boldsymbol{r})]=-M_S N_r m_{ir}\hat{r}=-M_S 2\pi d k_R m_{ir}$, where $N_r=2\pi d k_R$ and I have defined $H_d=4\pi M_S$. From Eq. (2) the two lowest branches of circularly symmetric coupled eigenmodes can be obtained. The eigenvalues obtained are in good agreement with the results obtained from dynamical micromagnetic simulations.[19] In particular, the frequency of the lowest mode, in which the radial components of the magnetization in the two layers are out of phase, is very accurately described by this simple model. The reason is that the most important approximations in this model are the shape of the magnetization density and that the assumed magnetization density is an eigenfunction of $\hat{h}_r$. In the lowest mode, the radial magnetostatic coupling is negligible since the radial components of the magnetization are out of phase. Therefore, the energy of this mode is controlled by the exchange bias and the interlayer exchange, which are well described by this simple model. Detailed comparison with extensive micromagnetic simulations[19] has verified that this simple model very accurately describes the behavior of the lowest-lying coupled eigenmode and that interactions not included in the model, such as interactions with the out-of-plane vortex cores, do not affect this mode in any significant way. The important term is the interlayer exchange interaction, which experiments show does depend on the thickness (and roughness) of the tunnel barrier. However, the thickness of the tunnel barrier varies only by a few Ångström for all relevant experimental ranges of interlayer exchange and product of area and resistance. Therefore, the tunnel barrier thickness is much smaller than thickness of the FM layers and than the radius of the disks, and the thickness of the tunnel barrier can be ignored.

I will use this model to study the effect of the spin torque arising from a spin-polarized current on the lowest eigenmode. In particular, I will focus on the behavior of the perpendicular spin torque as a function of bias voltage. With a bias voltage $V_b$ applied uniformly across the layers, a current spin-polarized current will flow, which will give rise to a torque on both layers. With the current taken as positive going from the PL to the FL, the torques on the PL and FL are,[4] respectively,

$$\boldsymbol{\tau}_{PL} = -|\gamma_e|\{a_J[\boldsymbol{m}_{PL}\times(\boldsymbol{m}_{PL}\times\boldsymbol{m}_{FL})] + b_J[\boldsymbol{m}_{PL}\times\boldsymbol{m}_{FL}]\},$$

$$\boldsymbol{\tau}_{FL} = |\gamma_e|\{a_J[\boldsymbol{m}_{FL}\times(\boldsymbol{m}_{FL}\times\boldsymbol{m}_{PL})] + b_J[\boldsymbol{m}_{FL}\times\boldsymbol{m}_{PL}]\}. \quad (3)$$

Here, $a_J$ and $b_J$ are the effective fields from the in-plane and perpendicular spin torques, respectively.

I first set the damping $\alpha$ to zero, as this simplifies the model while still correctly describing the behavior of the system. With zero damping we can also set $a_J=0$ since without damping $a_J$ only couples diagonally and does not give

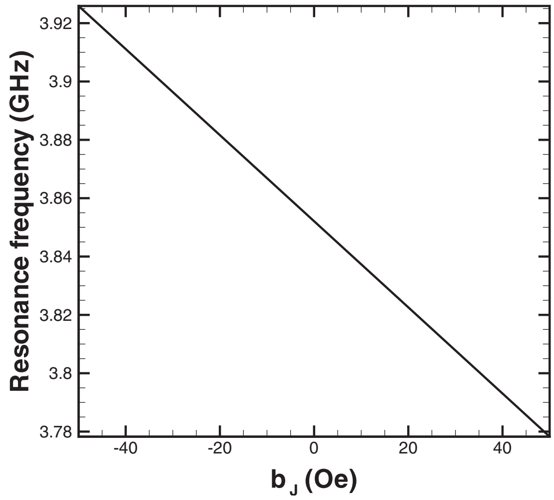


FIG. 1. Lowest resonance frequency as a function of perpendicular torque field $b_J$. In this figure, $R=1000$ nm, $d=12$ nm, $M_S=800$ emu/cm$^3$, $H_{IEC}=60$ Oe, and $H_{eb}=100$ Oe.

rise to a shift in frequency. By inserting these torques in the coupled LLG equations and linearizing, a simple secular equation for the two lowest eigenmodes is obtained

$$\begin{vmatrix} \gamma_e^2 A^2 - \omega^2 & \gamma_e^2 B^2 \\ \gamma_e^2 C^2 & \gamma_e^2 D^2 - \omega^2 \end{vmatrix} = 0,$$

where

$$\begin{aligned} A^2 &= [H_d + H_{eb} + H_{IEC} + b_J][M_S N_r + H_{eb} + H_{IEC} + b_J] \\ &\quad + [H_{IEC} + b_J][H_{IEC} - b_J - M_S Nr], \\ B^2 &= [H_d + H_{eb} + H_{IEC} + b_J][M_S N_r - H_{IEC} - b_J] \\ &\quad - [H_{IEC} + b_J][H_{IEC} - b_J + M_S Nr], \\ C^2 &= [H_d + H_{IEC} - b_J][M_S N_r - H_{IEC} + b_J] \\ &\quad - [H_{IEC} - b_J][M_S N_r + H_{eb} + H_{IEC} + b_J], \\ D^2 &= [H_d + H_{IEC} - b_J][M_S N_r + H_{IEC} - b_J] \\ &\quad + [H_{IEC} - b_J][H_{IEC} + b_J - M_S Nr]. \end{aligned} \tag{4}$$

In terms of these quantities, the lowest eigenfrequency is given by

$$\omega = |\gamma_e| \left[ \frac{A^2 + D^2}{2} - \frac{1}{2}\sqrt{(A^2 - D^2)^2 + 4B^2C^2} \right]^{1/2}. \tag{5}$$

By plotting $\omega$ as a function of $b_J$ for a range of reasonable values of $H_{IEC}$, $H_{eb}$, $M_S$, $d$, and $R$, it is clear that $\omega$ is a linear function of $b_J$ (see Fig. 1).

This means that one can determine $b_J$ experimentally for any given bias voltage and thus obtain the relationship between $b_J$ and bias voltage $V_b$. First, one measures the *zero-bias* resonance frequency $\omega_0$. The slope of $\omega$ as a function of $b_J$ can then be calculated directly

$$\beta \equiv \frac{d\omega}{db_J} = \frac{H_{eb}}{2\omega_0} - \frac{H_{eb} + H_d + 2H_{IEC}}{2\omega_0} \frac{A_0^2 - D_0^2 + B_0^2 - C_0^2}{\sqrt{(A_0^2 - D_0^2)^2 + 4B_0^2 C_0^2}}, \tag{6}$$

where $A_0$, $B_0$, $C_0$, and $D_0$ are the respective values at $b_J=0$. The parameters $H_{IEC}$, $H_{eb}$, $d$, and $R$ can all be determined independently. The frequency shift is given by $\Delta\omega = \omega - \omega_0 = \beta b_J$ and no adjustable parameters needed to determine the slope. By then measuring the eigenfrequencies at finite bias voltages, using the previously determined slope $\beta$ both magnitude and sign of $b_J$ can then be uniquely determined for each value of bias voltage (or bias current) by fitting the frequency shift $\Delta\omega$ to $\Delta\omega = \beta b_J$.

## III. MAGNETIC-SUSCEPTIBILITY MODEL

In a real experiment, the lowest-lying eigenfrequency can be measured through a dynamical susceptibility measurement in which an ac field is applied in the direction perpendicular to the disks and the response of the out-of-plane component of the magnetization density is measured. For enhanced signal, an array of disks can be patterned. By expanding the model used above to include finite damping as well as a time-dependent magnetic field $\boldsymbol{H}_a(t) = He^{-i\omega' t}\hat{z}$ with $\omega'$ the frequency of the applied field, the susceptibility, and the eigenmodes of the system can be calculated. With finite damping, both the real and imaginary parts of the eigenfrequencies of the susceptibility will shift slightly from the eigenfrequencies given by the simple model (5). The in-plane effective field $a_J$ will now contribute to a frequency shift but only of order $\alpha$. For small damping, the frequency shift is again linear in $b_J$ but its dependence on $a_J$ can be ignored. Similarly, the linewidth broadening depends on $a_J$ but its dependence on $b_J$ is negligible.

Specifically, in the presence of the applied field, the linearized LLG equation is now (suppressing the time dependence for ease of notation)

$$\frac{d}{dt}\begin{pmatrix} m_{1r} \\ m_{1z} \\ m_{2r} \\ m_{2z} \end{pmatrix} = \frac{|\gamma_e|}{1+\alpha^2}\begin{pmatrix} A_{r_1r_1} & A_{r_1z_1} & A_{r_1r_2} & A_{r_1z_2} \\ A_{z_1r_1} & A_{z_1z_1} & A_{z_1r_2} & A_{z_1z_2} \\ A_{r_2r_1} & A_{r_2z_1} & A_{r_2r_2} & A_{r_2z_2} \\ A_{z_2r_1} & A_{z_2z_1} & A_{z_2r_2} & A_{z_2z_2} \end{pmatrix}\begin{pmatrix} m_{1r} \\ m_{1z} \\ m_{2r} \\ m_{2z} \end{pmatrix} + \frac{|\gamma_e|}{1+\alpha^2}\begin{pmatrix} -H \\ \alpha H \\ -H \\ \alpha H \end{pmatrix}. \tag{7}$$

In the presence of spin torque, the matrix elements $A_{ij}$ above are given by

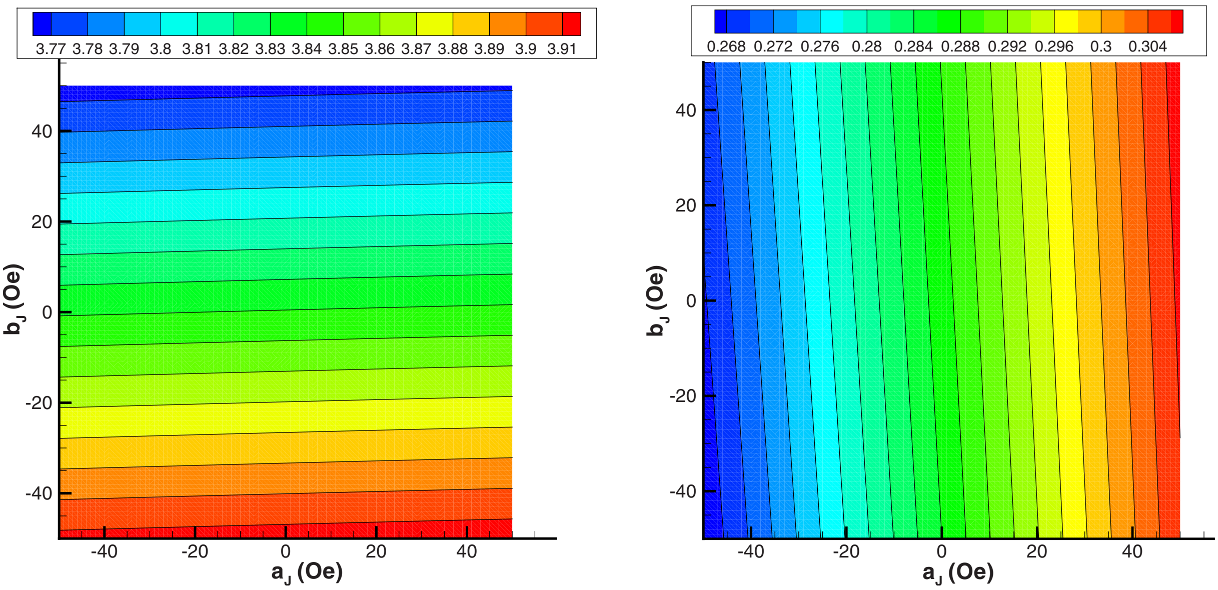


FIG. 2. (Color online) Real part (left panel) and imaginary part (right panel) of the lowest eigenmode frequency as functions of the in-plane and perpendicular torque fields $a_J$ and $b_J$. In this figure, $R=1000$ nm, $d=12$ nm, $M_S=800$ emu/cm$^3$, $H_{IEC}=60$ Oe, $H_{eb}=100$ Oe, and $\alpha=0.005$. The color coding indicates the frequency in gigahertz.

$$A_{r_1r_1} = -\alpha(H_{\mathrm{IEC}} - M_{\mathrm{S}}N_r + H_{\mathrm{eb}} + b_J) - a_J,$$

$$A_{r_1z_1} = H_d + H_{\mathrm{eb}} + H_{\mathrm{IEC}} + b_J - \alpha a_J,$$

$$A_{r_1r_2} = \alpha(H_{\mathrm{IEC}} + M_{\mathrm{S}}N_r + b_J) + a_J,$$

$$A_{r_1z_2} = -H_{\mathrm{IEC}} - b_J + \alpha a_J,$$

$$A_{z_1r_1} = -(M_{\mathrm{S}}N_r + H_{\mathrm{IEC}} + H_{\mathrm{eb}} + b_J) + \alpha a_J,$$

$$A_{z_1z_1} = -\alpha(H_{\mathrm{IEC}} - H_d + H_{\mathrm{eb}} + b_J) - a_J,$$

$$A_{z_1r_2} = -(M_{\mathrm{S}}N_r - H_{\mathrm{IEC}} - \mathrm{b_J}) - \alpha a_J,$$

$$A_{z_1z_2} = \alpha(H_{\mathrm{IEC}} + b_J) + a_J,$$

$$A_{r_2r_1} = \alpha(H_{\mathrm{IEC}} + M_{\mathrm{S}}N_r - b_J) - a_J,$$

$$A_{r_2z_1} = -H_{\mathrm{IEC}} + b_J - \alpha a_J,$$

$$A_{r_2r_1} = \alpha(M_{\mathrm{S}}N_r - H_{\mathrm{IEC}} + b_J) + a_J,$$

$$A_{r_2z_2} = H_d + H_{\mathrm{IEC}} - b_J + \alpha a_J,$$

$$A_{z_2r_1} = -(M_{\mathrm{S}}N_r - H_{\mathrm{IEC}} + b_J) + \alpha a_J,$$

$$A_{z_2z_1} = \alpha(H_{\mathrm{IEC}} - b_J) - a_J,$$

$$A_{z_2r_2} = -(M_{\mathrm{S}}N_r + H_{\mathrm{IEC}} - b_J) - \alpha a_J,$$

$$A_{z_2z_2} = -\alpha(H_{\mathrm{IEC}} - H_d - b_J) + a_J.$$

By solving this system of equations, the susceptibility, and the magnetization response can be obtained. Figure 2 depicts an example of the lowest eigenfrequency of the system as a function of $a_J$ and $b_J$. As is clear from the figure, the dependence of the real part (left panel) on $b_J$ is linear while the dependence on $a_J$ is negligible; conversely the dependence of the imaginary part (right panel) is linear in $a_J$ while its dependence on $b_J$ is negligible. This means that the simple model given by Eq. (5) is a very good approximation for the real part of the eigenfrequencies of the susceptibility.

Figure 3 shows an example of the real part of the susceptibility as a function of $a_J$ with $b_J=0.3a_J$. The frequency shift is very accurately given by $\Delta\omega=\beta b_J$, with $\beta$ given by Eq. (6). From the experimentally determined eigenfrequencies, the value of the perpendicular spin torque effective field $b_J$

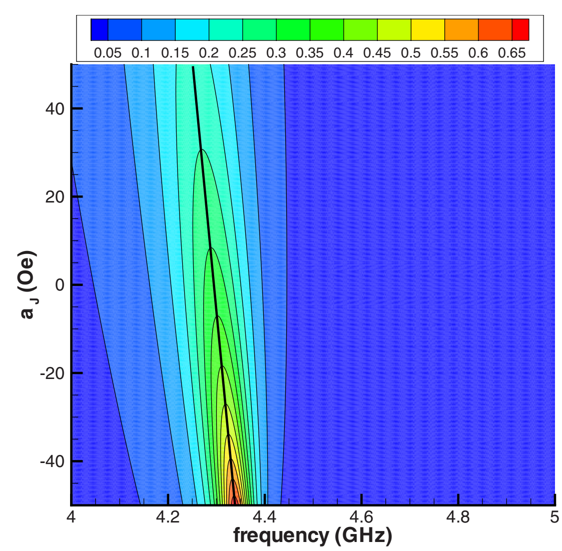


FIG. 3. (Color online) Contour plot of the real part of the perpendicular susceptibility for a system with $R=1000$ nm, $d=12$ nm, $M_S=800$ emu/cm$^3$, $H_{IEC}=60$ Oe, $H_{eb}=100$ Oe, and $\alpha=0.005$. The amplitude of the driving field is 100 Oe and here it was assumed that $b_J=0.3a_J$. The peak frequency is indicated by the black line, which is given by $\Delta\omega=\beta b_J$. The color coding indicates the amplitude of the susceptibility in arbitrary units.

can then be extracted and therefore its functional dependence on the bias voltage (or current) can be determined without any fitting parameters using the simple model Eq. (5). That is the main point of the present work.

## IV. DISCUSSION

Another experimental technique that can be used to determine $b_J$ as a function of bias voltage is to measure the signal noise spectrum of a single device under dc bias. The signal noise peak will not be at the eigenfrequency $\omega$, but at $2\omega$, because in the lowest eigenmode, the radial components of the PL and RL layers are almost perfectly out of phase while the $z$ components in each layer have a small phase difference from the radial components. This does give rise to a tunneling magnetoresistive signal but at twice the eigenfrequency. Note that if there were no phase difference between the $z$ and radial components in each layer, there would be no signal. The method of determining the effective field $b_J$ due to the perpendicular spin torque is the same. The measurements will give the eigenfrequency as a function of bias voltage. By independently calculating the slope $\beta$ of the eigenfrequency shift $\Delta\omega$ from the zero-bias eigenfrequency $\omega_0$ and using $\Delta\omega = \beta b_J$, the value of $b_J$ can be determined as a function of bias voltage. Additionally, the same set of parameters used to calculate $\beta$ can be used to check the value of $\omega_0$ checked for self-consistency between the measured value and that obtained from Eq. (5), both for the case of a susceptibility measurement as well as signal noise spectrum measurement.

A corresponding analysis can of course be carried out for an elliptical system. However, in order to make a comparison with a simple model, the system has to be assumed to be small enough that the magnetization is approximately uniform. This restricts the dimension of the major axis to be on the order of 200 nm or less, which requires electron-beam lithography or 193 nm UV lithography for good patterning control. A macrospin model can then be used but the demagnetizing field has to be determined in addition to other coupling parameters. Since the demagnetizing field depends sensitively on the shape of the system, it may be necessary to treat the demagnetizing field as an adjustable parameter, which makes direct comparison with experiments more complicated.

Finally, the change in eigenfrequency is not linear in $b_J$ as it is for the circularly exchange-biased disks, which makes the analysis and experimental extraction of $b_J$ vs $V_b$ more difficult for the elliptical system. The advantages of the system considered here are that it is a large and can easily be patterned with good control, the effects of the demagnetizing fields are well described by the simple model, and the coupling parameters can readily be determined independently. Furthermore, the analysis presented here is not based on a macrospin model but is a variational approach for the magnetization density across the disks.

It should be noted that I have here ignored both the effect of so-called Oersted fields induced by the current and spin torque acting on the out-of-plane vortex cores. For reasonable experimental systems, the Oersted fields are small and can be ignored. For example, for disk diameter $d = 1000$ nm, resistance-area product of 10 $\Omega\mu m^2$, and a bias voltage of 100 mV, the Oersted field is only 30 Oe at the edge of the system where the trial functions vanish. The size of the vortex cores is only on the order of 10 nm and the cores therefore occupy a very small fraction of the area of systems with radii of several hundred nanometers. As shown previously[19] the vortex cores do no affect the dynamics of the lowest eigenmode in the absence of spin torque—the spectra and eigenfrequencies are identical for systems with cores and systems in which the core region was removed. If the cores do not affect the eigenmodes in the absence of spin torque, the cores can safely be ignored in the presence of spin torque, which in the present work acts only as a perturbation on the eigenmodes.

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[1] J. C. Slonczewski, J. Magn. Magn. Mater. **159**, L1 (1996).
[2] L. Berger, Phys. Rev. B **54**, 9353 (1996).
[3] J. A. Katine, F. J. Albert, R. A. Buhrman, E. B. Myers, and D. C. Ralph, Phys. Rev. Lett. **84**, 3149 (2000).
[4] S. Zhang, P. M. Levy, and A. Fert, Phys. Rev. Lett. **88**, 236601 (2002).
[5] K. Xia, P. J. Kelly, G. E. W. Bauer, A. Brataas, and I. Turek, Phys. Rev. B **65**, 220401(R) (2002).
[6] M. A. Zimmler, B. Özyilmaz, W. Chen, A. D. Kent, J. Z. Sun, M. J. Rooks, and R. H. Koch, Phys. Rev. B **70**, 184438 (2004).
[7] S. Urazhdin, N. O. Birge, W. P. Pratt, and J. Bass, Phys. Rev. Lett. **91**, 146803 (2003).
[8] I. Theodonis, N. Kioussis, A. Kalitsov, M. Chshiev, and W. H. Butler, Phys. Rev. Lett. **97**, 237205 (2006).
[9] Z. Li and S. Zhang, Phys. Rev. B **69**, 134416 (2004).
[10] D. M. Apalkov and P. B. Visscher, Phys. Rev. B **72**, 180405(R) (2005).
[11] J. C. Sankey, Y.-T. Cui, J. Z. Sun, J. C. Slonczewski, R. A. Buhrman, and D. C. Ralp, Nat. Phys. **4**, 67 (2008).
[12] H. Kubota, A. Fukushima, K. Yakushiji, T. Nagahama, S. Yuasa, K. Ando, H. Maehara, Y. Nagamine, K. Tsunekawa, D. D. Djayaprawira, N. Watanabe, and Y. Suzuki, Nat. Phys. **4**, 37 (2008).
[13] S. Petit, C. Baraduc, C. Thirion, U. Ebels, Y. Liu, M. Li, P. Wang, and B. Dieny, Phys. Rev. Lett. **98**, 077203 (2007).
[14] C. Wang, Y.-T. Cui, J. Z. Sun, J. A. Katine, R. A. Buhrman, and D. C. Ralph, Phys. Rev. B **79**, 224416 (2009).
[15] C. Heiliger and M. D. Stiles, Phys. Rev. Lett. **100**, 186805 (2008).
[16] S.-C. Oh, S.-Y. Park, A. Manchon, M. Chshiev, J.-H. Han, H.-W. Lee, J.-E. Lee, K.-T. Nam, Y. Jo, Y.-C. Kong, B. Dieny, and

K.-J. Lee, Nat. Phys. **5**, 898 (2009).
[17] Z. Li, S. Zhang, Z. Diao, Y. Ding, X. Tang, D. M. Apalkov, Z. Yang, K. Kawabata, and Y. Huai, Phys. Rev. Lett. **100**, 246602 (2008).
[18] J. Sort, A. Hoffmann, S.-H. Chung, K. S. Buchanan, M. Grimsditch, M. D. Baró, B. Dieny, and J. Nogués, Phys. Rev. Lett. **95**, 067201 (2005).
[19] D. K. Schreiber, O. G. Heinonen, and A. K. Petford-Long, Phys. Rev. B **80**, 014411 (2009).
[20] M. Buess, T. P. J. Knowles, R. Höllinger, T. Haug, U. Krey, D. Weiss, D. Pescia, M. R. Scheinfein, and C. H. Back, Phys. Rev. B **71**, 104415 (2005).