Spin-transfer torques for domain wall motion in antiferromagnetically coupled ferrimagnets

Takaya Okuno1, Duck-Ho Kim1, Se-Hyeok Oh2, Se Kwon Kim3,4,*, Yuushou Hirata1, Tomoe Nishimura1, Woo Seung Ham1, Yasuhiro Futakawa5, Hiroki Yoshikawa5, Arata Tsukamoto5, Yaroslav Tserkovnyak3, Yoichi Shiota1, Takahiro Moriyama1, Kab-Jin Kim6, Kyung-Jin Lee2,7,8,* and Teruo Ono1,9,*

Antiferromagnetic materials offer ultrafast spin dynamics and could be used to build devices that are orders of magnitude faster than those based on ferromagnetic materials. Spin-transfer torque is key to the electrical control of spins and has been demonstrated in ferromagnetic spintronics. However, experimental exploration of spin-transfer torque in antiferromagnets remains limited, despite a number of theoretical studies. Here, we report an experimental examination of the effects of spin-transfer torque on the motion of domain walls in antiferromagnetically coupled ferrimagnets. Using a ferrimagnetic gadolinium–iron–cobalt (GdFeCo) alloy in which Gd and FeCo moments are coupled antiferromagnetically, we find that non-adiabatic spin-transfer torque acts like a staggered magnetic field, providing efficient control of the domain walls. We also show that the non-adiabaticity parameter of the spin-transfer torque is significantly larger than the Gilbert damping parameter, in contrast to the case of non-adiabatic spin-transfer torque in ferromagnets.

Antiferromagnets could provide a platform for spintronics with features distinct from their ferromagnetic counterparts1–3. In antiferromagnets, neighbouring spins are aligned antiparallel and therefore the magnetic dynamics and spin transport characteristics are expected to be fundamentally different from ferromagnets, where the spins are aligned parallel1. Recent experiments on magnetic dynamics have, in particular, shown that field-driven or spin–orbit-torque-driven4 domain wall (DW) motion in antiferromagnetically coupled ferrimagnets is fastest at the angular momentum compensation temperature (Tc), where the magnetic dynamics are antiferromagnetic. However, spin transport phenomena such as spin-transfer torque (STT) for antiferromagnetic DWs have so far only been explored in theoretical studies4–12.

Theoretical results suggest that STT for antiferromagnetic DWs consists of adiabatic and non-adiabatic STT components, as is the case for ferromagnetic DWs. In particular, theory suggests that non-adiabatic STT in antiferromagnets exerts a staggered magnetic field10,11; that is, it exerts effective magnetic fields of equal magnitude and opposite sign on the two different sublattices. Non-adiabatic STT for ferromagnetic DWs acts, in contrast, like a uniform magnetic field. The origin of non-adiabatic STT for antiferromagnetic DWs has been intensively debated not only theoretically but also experimentally over the past decade13–25. Nevertheless, the nature of non-adiabatic STT for antiferromagnetic DWs has not been investigated experimentally and remains an open question.

In this Article, we examine STT effects on DW motion in rare earth–transition metal (RE-TM) ferrimagnets across the Tc. We first explore DW motion in ferrimagnets theoretically, and conclude that non-adiabatic STT in antiferromagnets acts as a staggered magnetic field, which can drive the DWs. We experimentally confirm our theoretical results using a ferrimagnetic gadolinium–iron–cobalt (GdFeCo) alloy, in which Gd and FeCo moments are coupled antiferromagnetically, and find that STT effects can be used for efficient current-induced antiferromagnetic DW motion. We also find that the non-adiabaticity parameter (β) of STT is significantly larger than the Gilbert damping parameter (α), challenging conventional understanding of the non-adiabatic STT based on ferrimagnets.

Theoretical equation of ferrimagnetic DW velocity

We first describe the effect of STT on field-driven DW motion near the Tc of ferrimagnets based on the collective coordinate approach. The Landau–Lifshitz–Gilbert–Slonczewski-like equation for the magnetization of a ferrimagnet is given by26

\[ \dot{m} = \alpha m \times \dot{m} - \rho m \times \frac{\delta m}{\delta \sigma} \]

where m is the unit vector along the magnetization direction, \( \dot{m} \) is the equilibrium net spin density along –m, \( \alpha > 0 \) is the Gilbert damping constant, \( \sigma \) is the saturated spin density (the sum of the spin densities of two sublattices), \( \rho \) is the inertia associated with the dynamics of m, \( h_{eff} = -\delta U/\delta m \) is the effective field conjugate to m, and \( U[m] \) is the potential energy. The last two terms on the right-hand side are the adiabatic and non-adiabatic STT terms, where J is the charge–current density, P is the spin conversion factor given by \( P = (h/2e)(\sigma_1 - \sigma_2)/(\sigma_1 + \sigma_2) \) (with electron charge \( e > 0 \)), which

1Institute for Chemical Research, Kyoto University, Kyoto, Japan. 2Department of Nano-Semiconductor and Engineering, Korea University, Seoul, Republic of Korea. 3Department of Physics and Astronomy, University of California, Los Angeles, CA, USA. 4Department of Physics and Astronomy, University of Missouri, Columbia, MO, USA. 5College of Science and Technology, Nihon University, Funabashi, Chiba, Japan. 6Department of Physics, Korea Advanced Institute of Science and Technology, Daejeon, Republic of Korea. 7Department of Materials Science & Engineering, Korea University, Seoul, Republic of Korea. 8KU-KIST Graduate School of Converging Science and Technology, Korea University, Seoul, Republic of Korea. 9Center for Spintronics Research Network (CSRN), Graduate School of Engineering Science, Osaka University, Osaka, Japan. *e-mail: uzes.physics@gmail.com; kimsek@missouri.edu; kj_lee@korea.ac.kr; ono@scl.kyoto-u.ac.jp
represents the polarization of the spin-dependent conductivity \( \sigma_s (s=\uparrow\text{ or } \downarrow\text{ with } \uparrow\text{ chosen along } -\mathbf{m}) \) and thereby parameterizes the adiabatic torque term, and the non-adiabaticity \( \beta \) characterizes the non-adiabatic torque term. The DW velocity \( V \) above the Walker breakdown\(^{27,28}\) is obtained within the collective–coordinate approach (see Supplementary Note 1 for the derivation), given by

\[
V(\mu_H, J) = \frac{1}{\delta_T^2} \langle \alpha \lambda \mu_H H - \delta_T P J - \alpha \beta P J \rangle \tag{2}
\]

where \( M \) is the magnitude of the magnetization along \( \mathbf{m} \), \( \lambda \) is the DW width, \( \mu_H \) is the magnetic field and \( J \) is the magnitude of charge–current density. The first term is the DW velocity caused by the magnetic field and the second and third terms are the contributions of adiabatic and non-adiabatic STTs to the DW velocity, respectively. When a field and a current are applied simultaneously, one can separate pure field and STT effects by adding and subtracting the DW velocity \( V(\mu_H, J) \) for a positive current and the DW velocity \( V(\mu_H, -J) \) for a negative current, respectively:

\[
V_{\text{field}}(\mu_H) = \frac{V(\mu_H, J) + V(\mu_H, -J)}{2} = \frac{\alpha \lambda \mu_H H}{\delta_T^2} + \langle \alpha \lambda \rangle \tag{3}
\]

\[
\Delta V(J) = \frac{V(\mu_H, J) - V(\mu_H, -J)}{2} = -\frac{\delta_T P J}{\delta_T^2 + \langle \alpha \lambda \rangle} - \frac{\alpha \beta P J}{\delta_T^2 + \langle \alpha \lambda \rangle} \tag{4}
\]

Equation (4) shows that, at \( T_4 \) (that is, \( \delta_T = 0 \)), the adiabatic STT does not work for DW motion whereas the non-adiabatic STT acts like an effective field for the Néel order and thus moves the DW; this is consistent with previous theories for antiferromagnetic DWs\(^{8–12}\). We note that this effective field is staggered, meaning that the non-adiabatic STT induces the opposite effective magnetic fields on the two sublattices of antiferromagnets, and thus linearly couples to the staggered Néel order of antiferromagnetic DWs; in contrast, the non-adiabatic STT for ferromagnetic DWs is an effective uniform magnetic field that linearly couples to the local magnetization. Equation (4) also shows that the dependence of \( \Delta V \) on \( \delta_T \) is different for contributions of the adiabatic (first term) and non-adiabatic (second term) STTs. Assuming that the signs of both \( P \) and \( \beta \) do not change at \( T_4 \) (that is, \( \delta_T = 0 \)), which will be justified in the next section, the adiabatic contribution is antisymmetric with respect to \( \delta_T = 0 \) whereas the non-adiabatic contribution is symmetric. This allows us to separate two STT contributions by decomposing \( \Delta V \) into symmetric and antisymmetric components.

**Experiment on field-driven current-assisted DW motion**

To verify the above theoretical prediction, we investigated STT effects on current-assisted field-driven DW motion in ferrimagnetic GdFeCo compounds in which Gd and FeCo moments are coupled antiferromagnetically. Figure 1a presents a schematic illustration of our device. A 5 nm SiN/30 nm Gd\(_{40}\)Fe\(_{40}\)Co\(_{20}\)/100 nm SiN film was deposited on an intrinsic Si substrate by magnetron sputtering. The GdFeCo film was then patterned into 5-µm-wide and 500-µm-long microcires with a Hall cross structure using electron-beam lithography and Ar ion milling. For current injection, 100 nm Au/5 nm Ti electrodes were stacked on the wire. As the film lacks a non-magnetic heavy metal layer as a spin-current source, the effects of spin–orbit torque\(^{25,26,31}\) can be ignored. In this work, we focus on DW motion in the precessing regime, where the DW angle changes continuously (see Supplementary Note 2 and Supplementary Fig. 1 for a detailed discussion).

Figure 1b shows the DW velocity as a function of the magnetic field under positive and negative bias currents. We find that the DW velocity linearly increases with the magnetic field and is shifted under positive and negative bias currents. The results shown in Fig. 1b can be understood in terms of equations (2) to (4), that is, \( V(\mu_H, J) = V_{\text{field}}(\mu_H) + \Delta V(J) \). Following the analysis based on equations (3) and (4), \( V_{\text{field}}(\mu_H) \) and \( \Delta V(J) \) can be separated.

Figure 2a shows \( V_{\text{field}}(\mu_H = 85 \text{ mT}) \) as a function of temperature \( T \). It shows that \( V_{\text{field}} \) reaches its maximum at 241 K, irrespective of the current density \( |J| = 1.3 \times 10^5 \text{ A m}^{-2} \). The dotted green line indicates \( V_{\text{field}}(\mu_H) = (V(\mu_H, J) + V(\mu_H, -J))/2 \) and the orange arrow indicates \( 2\Delta V(J) = (V(\mu_H, J) - V(\mu_H, -J)) \).
density over the whole Fermi sea is not special for $p$, which is determined by the spin density only near the Fermi level. Therefore, $p$ and hence the spin conversion factor $P = \mu_B/2e$ cannot change signs at $T_A$. The assumption for $\beta$ can be understood similarly. The non-adiabaticity $\beta$ originates from the spin dissipation processes, which can be caused by several distinct microscopic Hamiltonian terms independent of the net spin density; examples include spin–orbit coupling and interactions with random magnetic impurities.\textsuperscript{10,19,22} This suggests that the sign of $\beta$ does not change at $T_A$.

To quantitatively compare the experimental results with the theory, we fitted the data in Fig. 2b using equation (4), where $\delta$ and $s$ were obtained by the measured $M$–$T$ curve (Supplementary Fig. 4).\textsuperscript{31} We considered $\alpha$, $\beta$ and $p$ as the fitting parameters, which are assumed to be temperature-independent for simplicity. The assumption of constant $\alpha$ is justified by our recent experimental report.\textsuperscript{5} Note that this is in contrast to an earlier experiment reporting that the effective $\alpha$ of ferrimagnets, estimated from the ferromagnetic resonance (FMR) linewidth, diverges at $T_A$ (ref. 34).

However, a recent theory\textsuperscript{35} suggests that the increase in FMR linewidth at $T_A$ can be attributed to the change in the nature of the magnetic dynamics from ferromagnetic (far away from $T_A$) to antiferromagnetic (at $T_A$), and not to the increase of $\alpha$ per se. The best fitting is the black solid line in Fig. 2b with the fitting parameters of $\alpha = (3.17 \pm 0.09) \times 10^{-1}$, $\beta = -0.53 \pm 0.02$ and $p = 0.109 \pm 0.002$. Figure 2c shows the adiabatic and non-adiabatic components in $\Delta V/J$ as a function of $T$, which are calculated from those parameters obtained by the fitting. The temperature dependence of the non-adiabatic component $\Delta V_{\text{STT}}/J$ is symmetric and that of the adiabatic component $\Delta V_{\text{ASTT}}/J$ is antisymmetric with respect to $T_A$, as expected from equation (4).

Mobility-based fitting-free analysis

Among the fitting values, a non-zero $p$ at $T_A$ is reasonable because it originates from the electron spins at the Fermi level, as explained above. On the other hand, we found a surprisingly large $|\beta|/\alpha$ of the order of 100, which is orders of magnitude larger than the values reported in most previous studies on ferromagnetic DWs.\textsuperscript{10,19,22} The adiabaticity $\alpha$ can be attributed to the change in the nature of the magnetic dynamics from ferromagnetic (far away from $T_A$) to antiferromagnetic (at $T_A$), and not to the increase of $\alpha$ per se. The best fitting is the black solid line in Fig. 2b with the fitting parameters of $\alpha = (3.17 \pm 0.09) \times 10^{-1}$, $\beta = -0.53 \pm 0.02$ and $p = 0.109 \pm 0.002$. Figure 2c shows the adiabatic and non-adiabatic components in $\Delta V/J$ as a function of $T$, which are calculated from those parameters obtained by the fitting. The temperature dependence of the non-adiabatic component $\Delta V_{\text{STT}}/J$ is symmetric and that of the adiabatic component $\Delta V_{\text{ASTT}}/J$ is antisymmetric with respect to $T_A$, as expected from equation (4).

![Fig. 2 | Field and current contributions to DW velocity as a function of temperature. a, $V_{\text{field}}(\mu_B H = 85 \text{ mT})$ (a) and $\Delta V/J$ (b) as a function of temperature under bias current densities of $|J| = 1.3, 1.7$ and $2.0 \times 10^{-9} \text{ A m}^{-2}$. Error bars for $\Delta V$ represent the standard deviation of $\Delta V$ values from a series of measurements under different magnetic fields (for example, seven $\Delta V$ values for the case of $|J| = 1.3 \times 10^{-9} \text{ A m}^{-2}$ and $T = 211.6 \text{ K}$, as shown in Fig. 1b). The fitting result of $\Delta V/J$ based on equation (4) is indicated by the black line. c, The non-adiabatic STT component $\Delta V_{\text{STT}}/J$ (red line) and the adiabatic STT component $\Delta V_{\text{ASTT}}/J$ (blue line) in $\Delta V/J$, calculated from the fitting result in b.](https://natur.elecrians/v2p0989/article/fig2.png)

\begin{equation}
\frac{\mu_C}{\mu_F} = \frac{\beta P}{\lambda M} - \frac{\delta J}{\alpha s} P
\end{equation}

The first and second terms correspond to the non-adiabatic and adiabatic STT contributions, respectively. At $T_A$ (that is, $\delta = 0$), the adiabatic STT contribution disappears while the non-adiabatic contribution exists. Figure 3a shows experimentally obtained $\mu_C/\mu_F$ as a function of temperature (black). To clearly show the finite $\mu_C/\mu_F$ at $T_A$, we extract the symmetric (red) and the antisymmetric (blue) components in $\mu_C/\mu_F$ as shown in Fig. 3a (Fig. 3b shows the symmetric component on a magnified scale). Although the data points are somewhat scattered, using the fact that $P$ and $\beta$ should vary smoothly with $T$ near $T_A$ without any singularity, as $T_A$ of ferrimagnets is not special for these spin transport parameters, we find that the symmetric component is obviously finite in the vicinity of $T_A$. We obtain $\mu_C/\mu_F(T \geq T_A)$ by averaging the data points within $T_A \pm 5 \text{ K}$, which yields $3.2 \times 10^{-14} \text{ m}^2 \text{T A}^{-1}$; thus, $\mu_C/\mu_F(T \geq T_A) = -\beta P/\lambda M = 3.2 \times 10^{-11} \text{ m}^2 \text{T A}^{-1}$. By substituting
of (ref. 34), we obtain symmetric (red) and antisymmetric (blue) components in as functions of temperature. Error bars for \( \mu \)

Spin transport in RE–TM ferrimagnets has been observed recently typical ferromagnetic DWs because of the antiferromagnetic align-

related to the underlying spin relaxation processes for ferromagnetic DWs, there has been longstanding debate about this \(|/\beta|\) also has an important physical meaning. For \( \beta \) velocity is proportional to \(|/\alpha|\)

Because the DW generates a DW because \( \mu_H \) is smaller than \( \mu_H \), the drive field does not reverse the magnetization or create DWs. Next, a current pulse (2 mA, 100 ns) was injected along the wire to generate an anomalous Hall voltage change. This value of \( \beta \) is very large, which can lead to fast current-induced antiferromagnetic DW motion, and the sign of \( \beta \) is negative, the origin of which needs further investigation. Our work shows that ferrimagnets can be useful for studying the magnetoelectronic properties of antiferromagnetically coupled systems, and also calls for further theoretical and experimental studies on STT for inhomogeneous antiferromagnetic spin textures.

Methods

Film preparation and device fabrication. The studied sample was an amorphous ferrimagnetic film of 5 nm SiN/30 nm GdFeCo–Co, 100 nm Si on an intrinsic Si substrate deposited by magnetron sputtering. To avoid oxidation of the GdFeCo layer, a 5-nm-thick SiN/Si layer was used as the buffer and capping layers. The film exhibits perpendicular magnetic anisotropy (PMA). GdFeCo microwires 5 µm wide and 500 µm long were fabricated using electron-beam lithography and an Ar ion milling process. A negative-tone electron-beam resist (maN-2403) mixed with a thinner for resists (T-1047, at a volume ratio of 1:1) was used for lithography at a wide and 500 µm for resists (T-1047, at a volume ratio of 1:1) was used for lithography at a fine resolution (~5 nm). For current injection, 100 nm Au/5 nm Ti electrodes were stacked on the wires.

Experimental set-up for field-driven DW motion. A pulse generator (Picosecond 10,300B) was used to generate the current pulse \( t_{\text{fix}} \) in Fig. 1a) and create the DW. A 2 mA, 100 ns current pulse was used to create the DW. For field-driven DW motion, a 2.5 or 3 mA bias d.c. current (with a corresponding current density of 1.3, 1.7 or 2.0 × 10^10 A m^-2) was injected along the wire to generate an anomalous Hall voltage. \( V_H \). The direction of the positive-bias d.c. current (+/− in Fig. 1a) is the same as that of the DW motion. A Yokogawa 7651 system that was used in a d.c. current source. The \( V_H \) at the Hall cross was recorded by an oscilloscope (Tektronix 7354) through a 46 dB differential amplifier. A low-temperature probe station was used to measure the DW motion for wide ranges of temperature.

DW detection technique. We used time-of-flight measurements of DW propagation to obtain the DW velocity in the flow regime. The procedure for measuring the DW velocity was as follows. First, a large perpendicular magnetic field \( H \) was applied to reset the magnetization. Next, a drive field \( H_C \) was applied in the direction of the DW motion, and \( H_C \) is the coercive field of the sample. Because \( \mu_H \) is smaller than \( \mu_H \), the drive field does not reverse the magnetization or create DWs. Next, a current pulse (2 mA, 100 ns) was injected by a pulse generator to create a DW next to the contact line through a current-induced Oersted field. We note that a 100 ns current pulse yielded a 100% DW writing probability for all temperatures and magnetic fields examined. As soon as the DW is created, the drive field \( \mu_H \) pushes the DW because \( \mu_H \) is larger than \( \mu_H \). Then the DW propagates along the wire and passes through the Hall cross region. When the DW passes through the Hall cross, the Hall voltage changes abruptly because the magnetization state of the Hall cross reverses as a result of the DW passage. This Hall signal change is recorded by the oscilloscope through a 46 dB differential amplifier. We refer to this as a ‘signal trace’. Because the detected Hall voltage change includes a large background signal, we subtract the background from the ‘signal trace’ by measuring a ‘reference trace’. The reference trace is obtained in the same manner as the signal trace except that the saturation field direction is reversed \( \mu_H = +150 mT \). In this reference trace, no DW is nucleated; hence, only electronic noise can be detected in the oscilloscope in the reference trace. The DW velocity \( v \) is determined as \( V = L/(t−t', \) where \( L \) is the DW travel length (400 µm in our measurement), \( t \) is the DW arrival time and \( t' \) is the DW generation time. \( t' \) was obtained as \( t' = (t_{\text{fix}} + t_{\text{sat}})/2, \) where \( t_{\text{fix}} \) and \( t_{\text{sat}} \) are

392 | VOL 2 | SEPTEMBER 2019 | 389–393 | www.nature.com/natureelectronics

M(T = \( T_0 \)) = \( 3.9 × 10^4 \) A m^-1 (Supplementary Note 5) and \( \lambda = 15 \) nm (ref. 11), we obtain \(-\beta = 0.058\). Here, the spin conversion factor \( P \) is replaced by the spin polarization \( p \). This value of \(-\beta \) is in good agreement with that obtained by the fitting, \(-\beta = 0.058 ± 0.003\), supporting the quantitative validity of the fitting values.

Both the fitting results and the mobility-based fitting-free analysis thus confirm the large \(|/\beta|/|/\alpha|\) quantitively. Because the DW velocity is proportional to \(|/\beta|/|/\alpha|\), this result suggests that highly efficient control of antiferromagnetic DWs by STTs can be achieved. The large ratio \(|/\beta|/|/\alpha|\) also has an important physical meaning. For ferromagnetic DWs, there has been longstanding debate about this ratio15–25, related to the underlying spin relaxation processes for \( \alpha \) and \( \beta \). We speculate that the observed large \(|/\beta|/|/\alpha|\) in our sample originates from the spin mistracking process with the small effective exchange interaction, which has been predicted to yield a large increase in \( \beta \) (refs. 17,18). In RE–TM ferrimagnets, the effective exchange averaged over two sublattices is smaller than that of typical ferromagnetic DWs because of the antiferromagnetic alignment of RE and TM moments. We note that a peculiar feature of spin transport in RE–TM ferrimagnets has been observed recently for spin–orbit-torque switching12 but not for DW motion. Another interesting observation is the negative value of \( \beta \) in GdFeCo, which is critically different from most18 known magnets. We speculate that this negative \( \beta \) may be related to the electron band structure of GdFeCo, as one theory23 has predicted that \( \beta \) can be negative in systems with both hole-like and electron-like carriers.

Conclusions

We have theoretically and experimentally examined the effects of STT on the motion of antiferromagnetic DWs in RE–TM ferrimagnets in the vicinity of \( T_x \). By controlling the temperature, it was possible to tune the net spin density in a GdFeCo ferrimagnet and thus distinguish between the adiabatic and non-adiabatic components of STT. Our experimental results show that the non-adiabatic STT in antiferromagnets acts like a staggered magnetic field, which can be used for efficient control of the antiferromagnetic DWs, confirming the theoretical prediction. Furthermore, we note two unusual properties of STT for the antiferromagnetic DWs in a ferrimagnet: the ratio of \(|/\beta|/|/\alpha|\) is very large, which can lead to fast current-induced antiferromagnetic DW motion, and the sign of \( \beta \) is negative, the origin of which needs further investigation. Our work shows that ferrimagnets can be useful for studying the magnetoelectronic properties of antiferromagnetically coupled systems, and also calls for further theoretical and experimental studies on STT for inhomogeneous antiferromagnetic spin textures.
the initial (final) time of the pulse current for DW generation. We averaged $V$ from five (when $230 < T < 250$ K) or three (otherwise) repeated measurements (note that $T_m = 244$ K in our sample). The error bar in $V$ was determined as the standard deviation of $V$ values from repeated measurements.

**Data availability**

All data that support the findings of this study are available from the corresponding authors on request.

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