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Effect of spin transfer torque on domain wall motion regimes in [Co/Ni] superlattice wires

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The combined effect of magnetic field and current on domain wall motion is investigated in epitaxial [Co/Ni] microwires. Both thermally activated and flow regimes are found to be strongly affected by current. All experimental data can be understood by taking into account both adiabatic and nonadiabatic components of the spin transfer torque, the parameters of which are extracted. In the precessional flow regime, it is shown that the domain wall can move in the electron flow direction against a strong applied field, as previously observed. In addition, for a large range of applied magnetic field and injected current, a stochastic domain wall displacement after each pulse is observed. Two-dimensional micromagnetic simulations, including some disorder, show a random fluctuation of the domain wall position that qualitatively matches the experimental results.

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The manipulation of magnetization by spin transfer torque (STT) [1,2] shows promise for the development of new data storage electronic devices. For instance, current-induced domain wall motion [3] may be implemented to increase the data storage efficiency. [Co/Ni] superlattices are often considered as a promising material for nanostructured spintronic devices because of their tunable magnetic and spin-electronic properties [5,8,9], especially for domain wall motion by STT [10–12]. However, in sputtered [Co/Ni] systems, specifically nanowires based on perpendicular anisotropy films, the current-induced DW motion has usually been studied in a restricted velocity regime. As a result, the adiabatic and nonadiabatic STT terms have been determined separately: Burrowes et al. [11] found that the nonadiabatic torque dominates the DW depinning and creep motion, whereas Koyama et al. [12] found the opposite, in the flow regime.

In this work, in order to extend the understanding of the effect of STT on DW motion, we propose a complete study in both velocity regimes, namely thermally activated and precessional flow. The sample used is a well-controlled [Co/Ni] superlattice stack consisting of Al2O3/Ni(0.5)/Au(1.2)/[Co(0.5)/Ni(0.6)]9/3/Au(1.2), with thicknesses in nanometers. The crystallinity and the layer-by-layer growth were monitored during deposition by recording reflection high-energy electron diffraction (RHEED) patterns and their intensity oscillation. Using a superconducting quantum interference device (SQUID) VSM magnetometer, the following parameters were determined: saturation magnetization $M_S = 9.3 \times 10^5$ A/m (930 emu/cm$^3$), effective anisotropy $K_{eff} = 3.0 \times 10^3$ J/m$^3$ (3.0 $\times 10^4$ erg/cm$^3$), and uniaxial magnetocrystalline anisotropy $K_U = 8.4 \times 10^3$ J/m$^3$ (8.4 $\times 10^4$ erg/cm$^3$) [14]. Gold was used as a capping and seed layer of the stack to avoid the spin-orbit torque observed in ferromagnetic nanowires with structural inversion asymmetry [20,21]. In order to study the current-induced DW motion, the sample was patterned into micron-wide wires (2–10 µm) by UV lithography and dry etching. An optical view of a 2-µm-wide wire is shown in Fig. 1(a). The 90-µm-long magnetic wire consists of two Hall crosses (not used in these experiments) and terminates in a triangular-shaped large pad for DW nucleation. Nonmagnetic contacts are patterned on each wire end for electrical current injection. The DW displacement is observed and quantified using magneto-optical Kerr effect (MOKE) microscopy. The magnetic configurations were imaged using a 25× magnification Leitz lens with 0.22 numerical aperture, which gives a nominal resolution around 1 µm. All experiments were done at room temperature.

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To quantify the current-induced DW motion, it is necessary to inject a high current density to observe a significant effect. The use of a microsecond-pulsed current ($I_{\text{pulse}}$) excitation alone is not enough to move the DW in the wire by STT; a pulsed magnetic field is also required. Using both field and current, the impact of STT on DW motion could be observed and analyzed. In our experiments, a homemade coil was used to apply microsecond-pulsed magnetic fields ($H_{\text{pulse}}$) homogeneously over the sample. In order to quantify the DW displacement, combined $H_{\text{pulse}}$ and $I_{\text{pulse}}$ excitations were synchronized. Due to the finite coil rise time, the $I_{\text{pulse}}$ trigger is delayed by 2 $\mu$s after the $H_{\text{pulse}}$ trigger. For all experiments, the synchronized pulse durations were 5 $\mu$s for $H_{\text{pulse}}$ and 3 $\mu$s for $I_{\text{pulse}}$, with a 2-$\mu$s delay [see Fig. 1(b)]. During the 2-$\mu$s magnetic field rise time, the applied magnetic field has a small impact on DW motion. The amplitude of each pulse was controlled using an oscilloscope [Fig. 1(b)].

The study of current- and field-induced DW motion was done on a single wire with 2 $\mu$m width. A sequence of DW displacements is shown in Fig. 1(c) for $J_{\text{pulse}} = +21 \times 10^{10}$ A/m$^2$ and $\mu_0 H_{\text{pulse}} = 20.5$ mT. The DW nucleation takes place on the right large pad, and the DW is injected in the wire by the right side using a low $H_{\text{pulse}}$. The upper picture in Fig. 1(c) shows the DW position at the entrance of the wire. Each subsequent frame shows, by image difference, the DW displacement after one single pulse. In this manner, we can measure the DW displacement after each pulse. The average DW velocity is calculated by dividing the displacement by the pulse duration (3 $\mu$s). When the DW displacement after a single pulse is too small to be observed, a series of pulses (up to 50) is used instead so as to observe a significant displacement. Since DW injection into the wire always occurs from the pad, by the current sign convention implies that a negative current is expected to favor by STT, field-driven DW propagation, whereas a positive current induces a STT which opposes the field-driven motion.

Measured DW velocities for various $H_{\text{pulse}}$ and $\pm J_{\text{pulse}}$ amplitudes are plotted in Fig. 2 in a semilogarithmic scale. Figure 2(a), resp. 2(b), shows negative, resp. positive $J_{\text{pulse}}$ (open symbols). Note that the curves for $J_{\text{pulse}} \geq 33 \times 10^{10}$ A/m$^2$ are in dotted lines because they correspond to a specific regime discussed at the end of the paper.

FIG. 1. (a) Optical picture of a typical 2-$\mu$m-wide wire studied here, with electrical connections. (b) Pulse synchronization for $H$ field (blue) and current (magenta): duration is 5 $\mu$s for $H_{\text{pulse}}$, 3 $\mu$s for $I_{\text{pulse}}$ with a delay of 2 $\mu$s. The current and field amplitudes can be tuned for each pulse with the following calibration: $V_{\text{PG}} = 14.6$ V amplitude for $\mu_0 H_{\text{pulse}}$ corresponds to 12.75 mT, and $V_{\text{PG}} = -15$ V for $I_{\text{pulse}}$ corresponds to a current of 9.68 mA. (c) Example of DW displacement sequence on a 2-$\mu$m wire width for $J_{\text{pulse}} = +21 \times 10^{10}$ A/m$^2$ and $\mu_0 H_{\text{pulse}} = 20.5$ mT. The black magnetic contrast corresponds to the DW displacement in the field direction during one pulse (difference of two pictures, taken at zero field before and after the pulse).

FIG. 2. Combined field and current-driven DW motion in a wire having a 2-$\mu$m width. (a) DW velocity as a function of $\mu_0 H_{\text{pulse}}$ in semilogarithmic scale for various (a) negative $J_{\text{pulse}}$ (full symbols) and (b) positive $J_{\text{pulse}}$ (open symbols). Note that the curves for $J_{\text{pulse}} \geq 33 \times 10^{10}$ A/m$^2$ are in dotted lines because they correspond to a specific regime discussed at the end of the paper.
mT [15]), possibly due to the effect of patterning. In Fig. 2, most of the data correspond to the transition between thermally activated and flow regimes [15], except in the high-positive-current case, where the magnetic field and the current seem to have no effect on the DW motion [24]. We will treat this unusual regime later separately. In the following we show the inset depicting the method for extracting the DW velocity at $\mu_0 H$ versus $J$ pulse in order to extract $\epsilon$ and $\eta$ parameters. The $\Delta H$ is a mean of the field shift for each velocity at $J_{\text{pulse}} \neq 0$. The black symbols are experimental data and the red line is a polynomial fit. The error bars are given by the standard deviation from the determination of $\Delta H$. Inset of (b): example of field shift determination for $\mu_0 H = 17.5 \text{ mT}$ for the velocity at $J = -9 \times 10^{10} \text{ A/m}^2$. (c) DW velocity versus $\mu_0 H_{\text{pulse}}$ for various $J_{\text{pulse}}$ in the precessional flow regime. (d) Plot of $v_J$ versus $J_{\text{pulse}}$, the inset depicting the method for extracting $\Delta v$, where the mean gives $2 \times v_J$ (for instance, $\Delta v$ is displayed for $J_{\text{pulse}} = +/− 9 \times 10^{10} \text{ A/m}^2$ at $\mu_0 H = 24 \text{ mT}$). The red line is a linear fit, and the error bars are the standard deviation from the determination of $v_J$.

To obtain a better understanding of how the thermally activated regime is affected by the injected current, the DW velocities only in this regime are plotted in Fig. 3(a) for both positive and negative $J_{\text{pulse}}$. For $J_{\text{pulse}} < 0$, the curves shift significantly downfield as compared to the zero-current curve (black crosses), i.e., for a given field, DW velocity increases with the amplitude of negative current. Thus, negative currents facilitate DW displacement, as expected for STT. On the other hand, for $J_{\text{pulse}} > 0$, no strong effect of the current is observed in Fig. 3(a). Moreover, whereas a decrease of DW velocity with respect to zero current would be expected for positive current, the DW velocity is larger for $J_{\text{pulse}} = +21 \times 10^{10} \text{ A/m}^2$. In order to quantify the effect of current in the thermally activated regime, we introduce an effective field $\Delta H$, the field shift with respect to a zero-current velocity curve (taken as reference). The field shift $\Delta H$ is measured as follows: for each $J_{\text{pulse}} \neq 0$, we measure a series of field shifts for varying values of $\mu_0 H$; the $\Delta H$ associated to that current density is their mean and the accuracy is their standard deviation. An example of field shift (red arrow) is displayed in the inset of Fig. 3(b) for the velocity at $J = -9 \times 10^{10} \text{ A/m}^2$ at $\mu_0 H = 17.5 \text{ mT}$. The field shift $\Delta H$ is counted positive for a shift to the left with respect to the zero-current velocity curve. In Fig. 3(b), we present the extracted $\Delta H$ values versus $J_{\text{pulse}}$ with the error bars. The experimental points can be fitted by a second-degree polynomial (red line): $\Delta H = +\epsilon J_{\text{pulse}} + \eta J_{\text{pulse}}^2$ with $\epsilon = -6.3 \pm 0.9 \times 10^{-15} \text{ T/A m}^2$ and $\eta = (3.3 \pm 0.3) \times 10^{-26} \text{ T/A m}^4$. It is equivalent to consider that the Arrhenius law can be rewritten as $\nu(H) = v_0 \exp[2M_S V_0(H_{\text{eff}} - H_{\text{dep}})/k_B T]$, where $H_{\text{eff}} = H_{\text{pulse}} + \epsilon J_{\text{pulse}} + \eta J_{\text{pulse}}^2$ is an effective out-of-plane field.

For the physical origin of such a $J$ dependence, we refer to Ref. [25], a theoretical analysis of the current-induced DW motion by magnetic field in the creep regime, including both nonadiabatic and adiabatic components of STT. A nonadiabatic term proportional to the current ($\epsilon J$) that acts as a magnetic field [26], and a quadratic adiabatic term ($\eta J^2$) which introduces nonlinearity to $H$, are added to the velocity expression in the creep regime. This theory was successfully used to interpret the experimental observation of the DW motion in GaMnAs nanowires [27,28]. From the...
efficiency $\varepsilon$, we can deduce $\beta$, the so-called nonadiabaticity STT constant, since $|\varepsilon| = (\beta PT)/(2eM_S\Delta)$ [26], where $P$ is the spin polarization ($P = 0.56$ being estimated in the flow regime, see below), and $\Delta = 5.8$ nm the DW thickness [29]. We obtain $\beta = (0.18 \pm 0.02)$, a value around 8 times bigger than reported on the sputtered [Co/Ni] system in the creep regime [11]. For sputtered [Co/Ni] systems, the relative contribution of both adiabatic and nonadiabatic terms is still unclear: Burrowes et al. [11] found that the nonadiabatic torque dominates the DW creep motion and depinning, whereas Koyama et al. [12] found that it plays no significant role. If we base our analysis on the theoretical expression from [25], we find that both adiabatic and nonadiabatic components may play a significant role on DW propagation in the thermally activated regime in epitaxial [Co/Ni] samples. The sign of the quadratic term $(\eta J^2)$ can be positive or negative according to Eq. (1) of Ref. [28]. The $J^2$ dependence could also relate to Joule heating. We use the following estimation of the temperature rise $\Delta T$ coming from Joule heating by a current pulse [30] $\Delta T = R I^2 T \times [\ln(16 K/((dC/w)^2)] + [\ln(t_{\text{pulse}})](2\pi K)$, which depends on $C, K$, and $d$, the specific heat, thermal conduction, and density of the substrate, respectively, the resistivity $\rho$ and dimensions ($l, w$) of the wire, and pulse duration $t_{\text{pulse}}$. Taking $d = 4000$ kg/m$^3$, $K = 40$ W m$^{-1}$ K$^{-1}$, $C = 700$ J kg$^{-1}$ K$^{-1}$ for sapphire and $R = 1500\Omega$, $w = 2\mu$m, and $l = 165\mu$m for our Co/Ni layer, one calculates that an injected current between 20 and $45 \times 10^{10}$ A/m$^2$ in our microwire gives rise to a slight temperature increase of about 4–20 K. The Joule heating contributes to an increase of temperature by $T + \Delta J^2$ with $\Delta$ constant. This affects the slope of the $v(J)$ curve in the thermally activated regime, because it decreases the potential energy barrier in the Arrhenius law. Therefore Joule heating does not cause the strong shift of the curves toward positive current that is observed.

As a first analysis of the data in the flow regime, we use the one-dimensional model with disorder proposed by Tatara et al. [18]. In this model, the DW configuration, moved by external field and current, is described by simply two coordinates: the DW position $X$ and the tilting angle $\Phi$ that the DW magnetization forms with the easy-axis plane. The following equations hold:

$$
\Phi' = \gamma H + \beta u/\Delta + f_{\text{pin}},
$$

$$
X' - \alpha \Delta \Phi = v_{\perp} \sin 2\Phi + u,
$$

where $H$ is the external magnetic field, $\alpha$ the damping parameter, $\beta$ the nonadiabaticity parameter of the STT, $\gamma$ the gyromagnetic ratio (1.8 $\times$ $10^{11}$ Hz/T for [Co/Ni]) [31]), $\Delta$ the DW width, $f_{\text{pin}}$ the pinning force and $v_{\perp}$ is a velocity related to the hard-axis magnetic anisotropy (also known as the Walker velocity). The last parameter $u = g P \mu_B J/(2eM_S)$ is a term proportional to the current density $J$, which is equivalent to a velocity (sometimes called the spin drift velocity). Here $g = 2.0$ is the Landé factor, $P$ the spin polarization, $\mu_B$ the Bohr magneton, and $M_S$ the saturation magnetization. The DW velocity at long times is the time average $v = \langle \dot{X} \rangle$ after depinning ($f_{\text{pin}} = 0$). Using the same method provided in the Supplementary Materials from Ref. [19], we deduce that the DW velocity can be expressed in the flow regime as the following sum:

$$
v = v_H + v_J,
$$

where $v_H$ and $v_J$ are, respectively, the field and current contribution given by

$$
v_H = \frac{\Delta \gamma H}{\alpha} \left( 1 - \frac{1}{1 + \alpha^2} \sqrt{1 - \left( \frac{H_W}{H} \right)^2} \right),
$$

$$
v_J = \frac{\beta u}{1 + \alpha^2} \left( 1 - \frac{\beta u}{\alpha} \sqrt{1 - \left( \frac{H_W}{H} \right)^2} \right),
$$

where $H_W$ is the Walker breakdown field. Equations (2), (3), and (4) allow the interpretation of our experimental data in the flow regime.

To get a better view of the impact of STT on the DW motion in the flow regime, the velocities in this regime are replotted in Fig. 3(c). For zero current, the velocity saturates at 13 m/s so $v_H$ is constant, as seen before in the case where the damping is very small [19]. Under field and current, we notice that the saturated velocity increases as the current density decreases, consistent with Eq. (2). Using this equation, we quantify $v_J$. Note that, for our Co/Ni system, we can simplify the expression [19]. Indeed, the damping parameter was measured to $\alpha \approx 0.02$ by ferromagnetic resonance (see also Ref. [31]), which leads to $\mu_0 H_w \approx 1.5$ mT. Thus, in the precessional flow regime, we have $\alpha^2 \ll 1$ and $H_W \ll H$, so that $v_J$ in Eq. (4) can be approximated to [19]

$$
v_J = u = \frac{g P \mu_B}{2eM_S} J.
$$

The processing of experimental data in the precessional flow regime in Fig. 3(c) should be made in the region where the velocity saturates (plateau). For $J_{\text{pulse}} < 9 \times 10^{10}$ A/m$^2$, this regime appears for strong fields ($\mu_0 H_{\text{pulse}} > 22$ mT), whereas outside this range, it appears at lower fields. To determine $v_J$ from the experimental data, we used the following method: the difference $\Delta \nu$ between saturated velocity for the same opposite current density should give $2 \times v_J$. For each opposite current, we measure a series of $\Delta \nu$ for various $\mu_0 H_{\text{pulse}}$, the mean giving $2 \times v_J$ and the standard deviation being the accuracy. The inset of Fig. 3(d) shows the method for extracting $\Delta \nu$ for $J_{\text{pulse}} = 9 \times 10^{10}$ A/m$^2$ at $\mu_0 H_{\text{pulse}} = 24$ mT.

The current contribution to the velocity $v_J$ is plotted as a function of current density in Fig. 3(d). We find that $v_J$ is proportional to $J$, in agreement with Eq. (5). The slope gives $v_J/J_{\text{pulse}} = (3.6 \pm 0.3) \times 10^{-11}$ m$^2$ A$^{-1}$ s$^{-1}$, corresponding to a spin polarization $P = 0.56 \pm 0.05$ (0.65 ± 0.1) was measured by spin-resolved photoemission spectroscopy [32] for the same stack). Using magnetotransport measurements, a spin polarization in the range 0.5–0.8 at 300 K (for instance see Refs. [19,33,34]) was reported in the literature for sputtered-growth [Co/Ni] multilayers due to thickness of both Co and Ni layers and stack number dependence [34]. These results confirm that the expressions of the above one-dimensional (1D) model are suitable for our system. Therefore, the experimental behavior of velocities observed in Fig. 2 can be explained in a large field range by a set of parameters found in the two independent regimes. This gives credit to our
that the structural symmetry of our system and the use of Au for the capping and seed layers are expected to suppress spin-orbit torques [20,21]. We note that due to sample design, Koyama et al. [19] could observe DW motion only in one direction of propagation and so could not detect the stochastic phenomena evidenced here.

This stochasticity extends over a large current density range, at least up to \( +4.5 \times 10^{10} \, \text{A/m}^2 \). Figure 4(c) shows all experimental displacements measured for \( J_{\text{pulse}} \geq +33 \times 10^{10} \, \text{A/m}^2 \). We have assigned positive values for DW displacement in the \( H \) field direction [Fig. 4(a)] and negative ones for DW displacement against it [Fig. 4(b)]. Zero displacements are also observed, meaning that sometimes the DW is pinned or the displacement is smaller than the optical resolution of 1 \( \mu \text{m} \). Small displacements can also take place by automation, where a structural change of the DW by STT leads to a DW displacement by itself [35], an effect also called DW inertia [36]. For [Co/Ni], using the DW width and damping quoted above, this automation is estimated around 1 \( \mu \text{m} \). The blue shaded regions displayed in Fig. 4(c) indicate the range where automation could occur and/or no DW motion by pinning.

The displacements displayed for various \( H_{\text{pulse}} \) and \( J_{\text{pulse}} \) in Fig. 4(c) form clouds, meaning that DW velocity is quasi-independent of \( J_{\text{pulse}} \). This result is unexpected according to the expression of \( v_J \) at long times. The statistical distribution of the set of DW displacements is shown in Fig. 4(d) without taking into account the value of the injected current density. This distribution is bell-shaped, with a slight skew asymmetry. The presence of a peak centered at zero is in favor of a random walk mechanism with compensation when \( v_H + v_J = 0 \). Note that the asymmetry may be a statistical artifact due to the small sample number of events (120 in total).

According to Eq. (5), \( v_J \) is proportional to \( J_{\text{pulse}} \), so that \( v_H + v_J = 0 \) should no longer hold for higher \( J_{\text{pulse}} \). For instance, taking \( J_{\text{pulse}} = +4.5 \times 10^{10} \, \text{A/m}^2 \), according to the slope \( v_J/J_{\text{pulse}} = (3.6 \pm 0.3) \times 10^{-11} \, \text{m}^3 \text{A}^{-1} \text{s}^{-1} \) deduced from Fig. 3(d), \( v_J \) should reach (16.2 \pm 1.4) m/s, leading to a velocity around \( v = (-7 \pm 1.5) \, \text{m/s} \). At this value, the DW should move against the field, not in the thermally activated regime. However, we measure [Fig. 2(b)] a DW velocity of around \( +2 \, \text{m/s} \). The DW motion is therefore slower than expected, which corresponds to the observed stochasticity and pinning. Thus, if we come back to the previous Fig. 2(b) depicting the complete \( v(H) \) curves for positives \( J_{\text{pulse}} \) (by keeping the \( H \)-driven case only), we attribute the strong change of the \( v(H) \) characteristic for \( J_{\text{pulse}} > +33 \times 10^{10} \, \text{A/m}^2 \) to the stochasticity. The decrease of the velocity for high current density cannot be explained using Eq. (5) unless we consider a huge drop of spin polarization (not consistent with the expected temperature dependence of \( P \) [33] due to the increase of temperature (max 20 K) coming from Joule heating produced by the current pulse). Thus, the 1D model is insufficient to describe the DW velocities in the stochastic regime. Investigation of current-induced DW motion in the nanosecond time scale [37] might reveal ultrafast back-and-forth movements of the DW.

To study the physics of this stochastic behavior, we have performed two-dimensional (2D) micromagnetic simulations of propagating DWs under field and current in a 1-\( \mu \text{m} \)-wide and 3.5-\( \mu \text{m} \)-thick track using the MUMAX3 code [38]. The
simulation parameters were as follows: cell lateral size 2 nm, damping factor 0.02, exchange stiffness 10 pJ/m, anisotropy constant $K_U = 8.4 \times 10^3$ J/m$^3$, and saturation magnetization $M_S = 9.3 \times 10^5$ A/m. To emulate the pinning induced by a variability in film thickness, we have divided the system into grains of random shape (mean diameter of 15 nm) of a different anisotropy parameter $K_U$ and saturation magnetization $M_S$ [see Fig. 5(b)]. For the disorder model stemming from an assumed variation of thickness, the values of $M_S$ and $1/K_U$ are assumed to follow a Gaussian distribution, with a standard deviation of 3% chosen so that the depinning field matches the experiments. For each set of magnetic field and current density values, we have repeated the simulation with different initial states: a homogenous Bloch wall pointing either up or down and “randomized” DWs containing sections pointing up or down with Bloch lines in between. The latter walls were prepared by submitting the homogeneous DWs to an intense easy-axis field for 1 ns. In all studied cases, the DW magnetization precessed nonuniformly along the DW length [shown in Fig. 5(c)], and all DWs quickly (<1 ns) evolved into nonhomogeneously magnetized DWs independently of the initial state. Homogeneous Bloch-up and Bloch-down DWs were tried in order to look for systematic automation-type displacements [35] in which, as stressed by Eq. (1), the DW velocity depends on the DW magnetization angle. As the experimental pulses are extremely long in comparison with the typical DW magnetization motion times (be it precession or even relaxation), such a dependence is not to be expected. However, in the presence of disorder and for the long DWs considered, that are beyond the 1D physics leading to Eq. (1), this has to be numerically tested.

We observe that the progression of the DW position is noisy as a consequence of the precessions induced by the random variation of material parameters. To illustrate this, we have repeated 30 simulations with different random grains and three different DW initial states (Bloch wall pointing up, down, and randomized), for $J = 2.5 \times 10^{11}$ A/m$^2$ and $\mu_0 H_z = 25$ mT. Figure 5(a) shows the typical dispersion of DW position with time. Although the DW velocity varies randomly, we observe that the mean velocity is linearly dependent on the applied current. In the inset of Fig. 5(a), we show the progression of a DW for different current densities (under $\mu_0 H_z = 25$ mT). The random fluctuation of the DW position qualitatively matches the experimental results shown above. Quantitatively, however, the DW position fluctuations are smaller than seen in experiments and, more importantly, the zero average DW displacement is obtained only in a restricted vicinity of the current, when it compensates the applied field effect. It is possible that this discrepancy simply stems from the simplicity of the disorder model used here, but this requires a systematic simulation study that is beyond the scope of this experimental paper.

To conclude, we have experimentally investigated the combined effect of field and current on DW motion in epitaxial [Co/Ni] samples with strong perpendicular anisotropy, in microwire devices and using Kerr microscopy. A strong impact of both amplitude and polarity of current on DW propagation, in two different velocity regimes, has been observed. The experimental data may be explained by the presence of both nonadiabatic and adiabatic terms in the STT expression. In the flow regime, the behavior of DW velocity is consistent with the 1D model. At high currents, when opposing the field effect, a stochastic behavior is observed. Real-scale 2D micromagnetic simulations show a random fluctuation of the DW position that qualitatively matches the experimental results. Surprisingly, this stochasticity experimentally extends over a large field and current range.

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FIG. 5. Micromagnetic simulations of DW propagation in an inhomogeneous film. (a) DW progression for 30 cases with different random grains and different DW initial states, for $\mu_0 H_z = 25$ mT and $J = 2.5 \times 10^{11}$ A/m$^2$. The colors correspond to the initial DW states (blue for Bloch-up, green for Bloch-down, and yellow for randomized DW). The inset shows the progression for different current densities (blue for $\mu_0 H_z = 25$ mT). (b) Map of the random grains (the gray scale corresponds to the magnitude of $M_s$). (c) Three snapshots of the magnetization corresponding to one of the curves in (a).
EFFECT OF SPIN TRANSFER TORQUE ON DOMAIN WALL . . .


[29] DW thickness is estimated as : with \( A = 10 \times 10^{-12} \text{J/m} \) the exchange constant for Co/Ni.


