

Effect of Ir growth pressure on the domain wall dynamics in Ta/Pt/Co/Ir/Ta stacks

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The dynamical response of magnetic domain walls in ultrathin magnetic films to external magnetic fields is determined not only by the composition and thickness of the layers but also by the growth conditions. Growth conditions can induce significant structural changes to the layers and at the interfaces between them, affecting in particular the dynamics of domain walls, their mobility, elastic tension, and the pinning forces acting on them. In this work, we focus specifically on the effect of Ir layer growth pressure in Ta/Pt/Co/Ir/Ta ultrathin multilayered films. Measurements of the DC magnetic properties, domain-wall velocity, and domain morphology in the creep regime for both constant and alternating field pulses were performed for a batch of samples where the Ir layer was grown at different pressures. We find that the saturation magnetization, the effective anisotropy constant, and the domain-wall surface tension grow with increasing pressure and saturate at a threshold pressure, while the Dzyaloshinskii-Moriya field and the strength of the disorder remain practically unaltered over the range of pressures considered.

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I. INTRODUCTION

The special properties that emerge when one or more characteristic dimensions of a magnetic system are reduced to the nanometer scale have attracted significant attention in recent years due to their promising technological applications [1]. Among the most studied examples of such systems are ultrathin ferromagnetic films, which consist of multilayer stacks with layers made of different materials of nanometric thicknesses. These systems are particularly sensitive to the orientation and strength of magnetization [2–6], and thus potentially attractive for novel devices. However, their development entails complex nanofabrication methods. Both the film composition and growth conditions have been shown to strongly influence the magnetic properties, type of domain wall (DW) stabilized, and DW propagation [7–9]. Strikingly, despite their sensitivity, magnetic thin films with perpendicular anisotropy have also become a model experimental system for studying the universal dynamical properties of one-dimensional extended elastic interfaces in random media [10]. This universality provides a convenient phenomenological framework for exploring the effect of growth conditions on DW dynamics by focusing the attention on the constitutive equations of a few experimentally accessible transport coefficients.

The dynamics of magnetic DWs are mainly determined by the interplay between elasticity, external fields, thermal fluctuations, and quenched disorder. The surface tension and the type and strength of the pinning forces acting on a

domain wall critically depend on microscopic properties, including bulk and surface anisotropies, asymmetries, spatial heterogeneities and defects, dipolar interactions, exchange interactions, and the Dzyaloshinskii-Moriya interaction (DMI). The presence of DMI in ultrathin films broadens the possibilities of these systems further, owing to its relevance in generating intriguing magnetic textures such as chiral domain walls and skyrmions [11]. These small magnetic textures hold significant potential for application in the next generation of digital processing and recording techniques based on spintronics, which aim to reduce power consumption while enhancing capacity, processing speed, and robustness [2].

Among the film stacks that have been studied are Pt/Co films with top layers composed of heavy metals (HMs) such as Pt, Ni, Ta, Ir, among others [11–17]. The effects of changing Ir thickness in Ta/Pt/Co/Ir/Ta stacks on the resulting DMI and domain-wall dynamics were recently analyzed in Ref. [17], revealing substantial changes in DMI strength. Another interesting growth modification is to use different Ar pressures when growing the top layer of the system. There are few studies that demonstrate a significant change in the dynamic properties of the samples [12,16], but particularly none in Pt/Co/Ir films. In all of these works the magnetic characteristics of the samples are analyzed in the presence of a constant magnetic field (DC field), but the response to an alternating magnetic field (AC) was not explored. Recently, it was shown that the AC response yields valuable information about domain walls and pinning properties. These properties

are complementary to those obtained by using DC fields, as demonstrated in Refs. [18,19].

In this paper, samples of $\text{SiO}_2/\text{Ta}(5)/\text{Pt}(3)/\text{Co}(0.8)/\text{Ir}(1)/\text{Ta}(2)$ (thicknesses in nm), are investigated. Samples were grown via DC magnetron sputtering with different Ar pressures used during the growth of the Ir layer; we refer to the Ir growth pressure as p_{Ir} . All samples have the same nominal thickness for each layer. We analyze the changes observed in the hysteresis loops, static micromagnetic properties, domain-wall velocities under a DC or AC field, and DMI intensity through asymmetric domain growth in presence of in-plane and out-of-plane fields.

II. PROPERTIES OF INTEREST

In perpendicularly magnetized ultrathin films with bubble-like domains, the study of the domain-wall (DW) velocity as a function of a constant and uniform magnetic out of plane field H is a standard probe to study dynamical properties of DWs. The velocity-field characteristics $V(H)$ thus obtained reveals several dynamical regimes with interesting universal features [20–22]. These universal features manifest at low velocities, in the creep [21] and depinning [22] regimes, for $H \ll H_d$ and $H \sim H_d$, respectively, with H_d the depinning field. The most robust universal property is unveiled in the low-field regime with $T > 0$ and $H \ll H_d$ which is well characterized by the creep law

$$V(H) = V_0 \exp[-\alpha(\mu_0 H)^{-\mu} + \beta], \quad (1)$$

where V_0 is a characteristic velocity, α and β are material- and temperature-dependent but field-independent parameters and μ_0 the vacuum permeability. This law describes a thermally activated motion of thermal nuclei over typical energy barriers $U_{\text{opt}} \sim \alpha T(\mu_0 H)^{-\mu}$. The creep exponent is predicted to be $\mu = 1/4$ for one-dimensional interfaces in random-media with the so-called “random-bond” type of disorder [23–26]. The exponent $\mu = 1/4$ in particular is consistent with a large number of measurements performed in ultrathin films of different materials [8,11,17,27–32]. These experiments confirm the predictions of elastic interface models, the effective dimensionality of the DWs and the type of disorder present in the samples. Indeed, a “random-field” type of disorder, a nonlocal elasticity or a different DW dimensionality is also well predicted by Eq. (1) but with a clearly distinguishable $\mu \neq 1/4$ [26]. The creep formula of Eq. (1) thus offers a robust phenomenological framework to analyze the effect of growth conditions in the pinning and elastic properties of DWs. Furthermore, α and β can be predicted by the Larkin collective pinning theory [24,33,34]. Using Larkin collective pinning theory [24,33] it is possible to connect α and β with different parameters such as the DW surface tension σ , the saturation magnetization M_s , the DW width δ as well as with relevant pinning parameters such as the strength of the disorder Δ_0 and the pinning correlation length ξ (we describe and exploit these relations in Sec. V). All these parameters are in principle sensitive to growth conditions and in particular p_{Ir} . An example of the effect of growth pressure on Pt/Co/Pt samples is reported by Lavrijsen *et al.* [12], where growth conditions are shown to strongly impact the DW velocity profiles

of expanding bubbles under the simultaneous application of out-of-plane and in-plane magnetic fields.

In perpendicularly magnetized systems, by applying an in-plane magnetic field (H_x) in addition to an out-of-plane field H , asymmetric growth of bubble domains can be observed if DMI is present [13,35–37]. In these systems, DWs grow with different speeds depending on their orientation relative to H_x . When velocity is measured parallel to H_x a minimum is observed when $H_x = H_{\text{DMI}}$,

$$\min_{H_x} V(H, H_x) = V(H, H_{\text{DMI}}), \quad (2)$$

with $V(H, H_x)$ being the velocity-field characteristics under the tilted field, such that $V(H, H_x = 0) \equiv V(H)$ of Eq. (1) in the creep regime. From H_{DMI} we can estimate the DMI strength as

$$D = \mu_0 H_{\text{DMI}} M_s \delta, \quad (3)$$

with $\delta = \sqrt{A/K_{\text{eff}}}$ the DW width, A the exchange interaction, K_{eff} the effective anisotropy and M_s the saturation magnetization [38]. An isotropic growth implies $H_{\text{DMI}} = 0$, and thus, a negligible DMI. In our case, the measurement of H_{DMI} for samples with different p_{Ir} allows us to study not only the role of the interface engineering on the strength of the DMI but also to test theoretical predictions relating it to changes in the total surface tension σ , which we can experimentally estimate as explained in the following.

Complementary information on the pinning and elasticity of DWs can be obtained by applying a square wave of an out of plane magnetic field $H(t)$ to a compact domain [18]. In the absence of any magnetic field, a micrometer sized domain in an ultrathin disordered film is typically trapped into a metastable state from which it cannot escape within experimental timescales, in spite of having a negative average curvature promoting its collapse. The applied pulses are hence chosen to be large enough to induce an AC-assisted curvature driven collapse of the DW in experimental timescales. This method exploits the approximation that the local response of a DW to an external field is given approximately by its DC velocity response $V(H_{\text{eff}})$ in an effective local field $H_{\text{eff}} = H(t) + C\kappa$, where κ is the local signed curvature of the DW and $C = \sigma_C/2M_s$, with M_s the saturation magnetization and σ_C a phenomenological elastic tension, expected to be of the same order of magnitude than the micromagnetic prediction $\sigma_0 \approx 4\sqrt{AK_{\text{eff}}}$, with A the exchange interaction and K_{eff} the effective anisotropy coefficients. Under this approximation and also assuming that the area oscillations are small compared with the total area it can be shown that at time $t = N\tau$ we have the effective area [19]

$$\Lambda(N) \equiv -\frac{\mathcal{A}(N) + \mathcal{P}(N)V(H)(\tau/4)}{2\pi V'(H)\tau} \approx \Lambda(0) + CN, \quad (4)$$

where $\mathcal{A}(N)$ is the average area of the domain, $\mathcal{P}(N)$ is the average perimeter of the DW at the N th pulse, τ the period of the square wave, $V(H)$ is the DC velocity-field characteristics, and $V'(H)$ its derivative with respect to the AC field amplitude H . Equation (4) generalizes the exact and universal spontaneous decay $A_t = A_0 - 2\pi Cmt$ expected for the area of any compact domain of initial area A_0 when the velocity response is simply $V(H) = mH$, with m the DW mobility, a situation

typically realized in the absence of pinning at relatively small or large fields. Numerical simulations of an elastic DW in presence of an AC field show that in the absence of disorder Eq. (4) is very accurate. In the presence of disorder however, it holds only for small N , provided the initial condition is not strongly correlated with the underlying disorder. The later condition is typically realized just after nucleation of a compact domain by a short and high magnetic field pulse applied in a saturated magnetization background. We can thus use Eq. (4) to extract C by only monitoring $\mathcal{A}(N)$ and $\mathcal{P}(N)$, both accessible from images. Furthermore, as long as we have access to M_S , we can compute σ_C .

For larger N , the prediction $\Lambda(N) - \Lambda(0) \sim CN$ is expected to fail due to the roughening of the DW. For nearly circular initial domains, DW roughening can be monitored in the same experiment by using the quadratic mean displacements of the DW with respect to a perfect circle with a radius equal to the angularly averaged time-dependent radius of the deformed DW loop. If $r(\theta, N)$ is a univalued function that describes the DW in polar coordinates from the nucleation center in the N th pulse, we can define the N -dependent roughness as

$$\langle u^2 \rangle = \langle r^2 \rangle_\theta - \langle r \rangle_\theta^2, \quad (5)$$

where $\langle \mathcal{F} \rangle_\theta \equiv (2\pi)^{-1} \int_0^{2\pi} \mathcal{F} d\theta$. The quantity $\langle u^2 \rangle$ is expected to grow with N and to scale as a generic roughening process $\langle u^2 \rangle \sim N^\gamma$, with γ being a growth exponent [39].

To complete the set of properties just described we also analyze the saturation magnetization M_S and the anisotropy field H_K . All the above properties are measured as a function of p_{Ir} .

III. EXPERIMENTAL DETAILS

We studied ultrathin films with perpendicular magnetic anisotropy (PMA) grown by DC magnetron sputtering on thermally oxidized SiO_2 substrates. The structure of the samples was $\text{SiO}_2/\text{Ta}(5)/\text{Pt}(3)/\text{Co}(0.8)/\text{Ir}(1)/\text{Ta}(2)$ (the nominal thickness of each layer is specified in nm in parentheses). All layers were grown using Ar pressure of 3 mTorr, except for the Ir layer, which was grown with different pressures p_{Ir} , ranging from 3 to 25.5 mTorr while maintaining a constant nominal thickness for the Ir layer. All layers were grown using a sputtering power of 50 W. Before deposition, the sputtering rate of each material was calculated by growing a film for ten minutes and performing x-ray reflectivity (XRR) measurements to determine the thickness of the deposited material. A separate growth rate was calculated for each p_{Ir} employed. By performing XRR and rocking curve measurements of the multilayer structures, the total thickness of the films and the roughness and elemental intermixing were assessed.

The magnetic properties of each sample were characterized using two techniques: A superconducting quantum interference device (SQUID) to measure the static magnetic properties, and a homemade polar magneto-optical Kerr effect (PMOKE) microscope to image the magnetic domain morphology of the samples and characterize the dynamics and morphology of domain walls. Hysteresis loops as a function of the in-plane and out-of-plane magnetic field were measured using SQUID at room temperature. This allowed us to obtain the saturation magnetization values M_S , the effective

anisotropy field H_K to calculate the effective anisotropy constant $K_{\text{eff}} = \mu_0 H_K M_S / 2$ [40] and confirm that all the samples presented a dominant PMA. Furthermore, to analyze the morphology of the domains during reversal, images were taken using PMOKE microscopy. To build the PMOKE hysteresis loop we applied 50 ms square pulses of increasing and decreasing amplitudes in steps of $\mu_0 \Delta H = 0.5$ mT, and images were taken between pulses.

The DW dynamics were characterized by PMOKE using three different and complementary approaches. All protocols start by nucleating a compact magnetic domain by applying out-of-plane magnetic field pulses of magnitudes on the order of 20 mT, with a duration of 5 ms, opposed in direction to a previously out-of-plane magnetically saturated background.

(i) In the first DC characterization, we apply a sequence of out-of-plane magnetic field square pulses of amplitude H and different durations to grow the nucleated domain in successive steps [see, for instance, the inset of Fig. 3(b)]. As described in Ref. [18], the sequence of average DW displacements for a given pulse duration is used to estimate the average velocity $V(H)$ for $4 \text{ mT} \leq \mu_0 H \leq 35 \text{ mT}$, and in particular to determine the constants α and β associated with the creep regime through Eq. (1).

(ii) In the second DC characterization, we apply a sequence of identical out-of-plane magnetic field square pulses of amplitude H to grow the nucleated domain in steps, in the presence of a constant in-plane magnetic field H_x . We thus obtain $V(H, H_x)$, for $-300 \text{ mT} \leq \mu_0 H_x \leq 300 \text{ mT}$ and $\mu_0 H = 100 \text{ mT}$, and H_{DMI} determined using Eq. (2).

(iii) In the third AC characterization, a pure (i.e., without a DC component) out-of-plane square-wave magnetic field is applied. Every period of the square wave, of duration $\tau = 50$ ms, consists of a positive square pulse of duration of 25 ms immediately followed by a negative square pulse of the same amplitude and duration. The amplitude of the square wave is fixed so to always produce a fixed $6 \mu\text{m}$ displacement in the first half-period positive pulse. Since the curvature of the DW plays a relevant role in this protocol, we ensure that the initially nucleated domain has a fixed $\approx 40 \mu\text{m}$ radius, and then acquire images every period to determine the area \mathcal{A} and perimeter \mathcal{P} of the domain, and then the Λ function from Eq. (4) as a function of the period number N to finally estimate C .

IV. RESULTS

A. Structural characterization

The low-angle XRR data for samples with $p_{\text{Ir}} = 3$ and 25.5 mTorr can be seen in Fig. 1(a). The periodicity of the oscillations (Kiessig fringes) is sensitive to the total thickness of the multilayer (t_S) [41]. The value of t_S can be obtained fitting the modified Bragg's law [27,42]:

$$\sin(\theta_n)^2 = 2\delta + n^2 \left(\frac{\lambda}{2t_S} \right)^2, \quad (6)$$

where λ is x-ray wavelength, $1 - \delta$ is the real part of the refractive index of the film, and n is the order of each local minima, θ_n [see inset in Fig. 1(a)]. From this technique a total

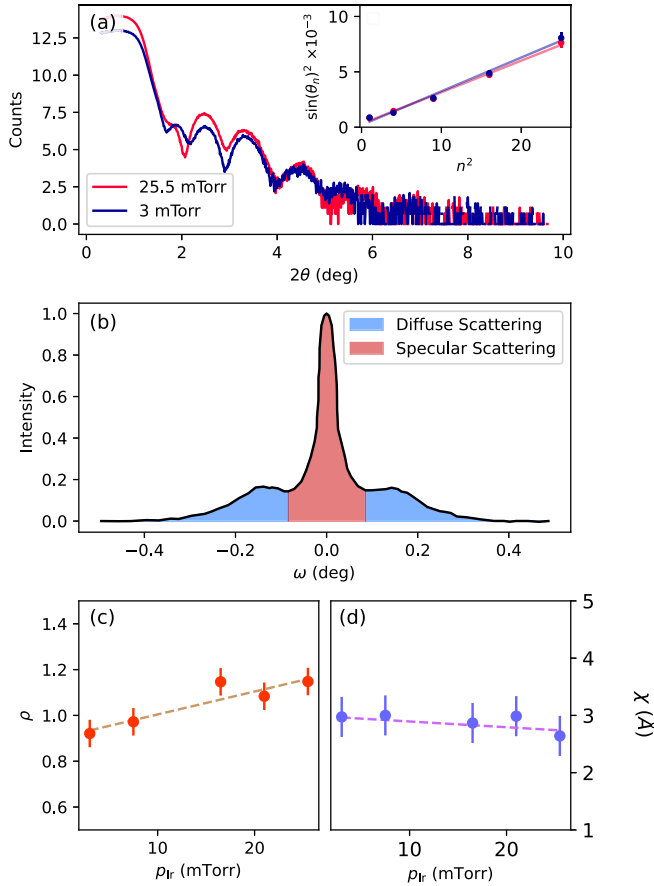


FIG. 1. (a) Low-angle x-ray reflectivity data for Pt/Co/Ir structures deposited with $p_{\text{Ir}} = 3$ and 25.5 mTorr. The inset shows the linear fits using Eq. (6) and the obtained values for t_s . (b) Rocking curve for $p_{\text{Ir}} = 3$ mTorr, acquired around the first Kiessig fringe. Different colors corresponds to diffuse and specular scattering. Global roughness parameter ρ and χ factor, both as a function of p_{Ir} are shown in panels (c) and (d), respectively.

value of the layer growth on the SiO_2 substrate of $t_s = 10 \pm 1$ nm was obtained.

In Fig. 1(b), a typical rocking curve measurement is shown for the sample grown with $p_{\text{Ir}} = 3$ mTorr. This rocking curve was acquired around the first Kiessig fringe determined from the XRR measurements. From the rocking curves, the areas corresponding to diffuse scattering (S_d) and specular scattering (S_s) were calculated. The ratio $\rho = S_d/S_s$ is proportional to a global roughness of the multilayer [43]. The ρ values as a function of p_{Ir} are shown in Fig. 1(c). From the figure, an increase in the roughness (ρ) of the interfaces as p_{Ir} increases is evinced. This trend may be attributed to the increase of the Ir grain size with increasing deposition pressures due to the lower kinetic energy of the Ir atoms arriving at the substrate [44]. To further quantify the interfacial quality of the Pt/Co/Ir structures, the XRR data was analyzed using GenX [45], from which a roughness-intermixing parameter χ can be calculated.

In thin-film systems, roughness typically describes the height of variations in the structural morphology at an interface. In contrast, intermixing describes the thickness to which chemical species are alloyed at an interface. It is important to note that the χ parameter calculated from refinement of

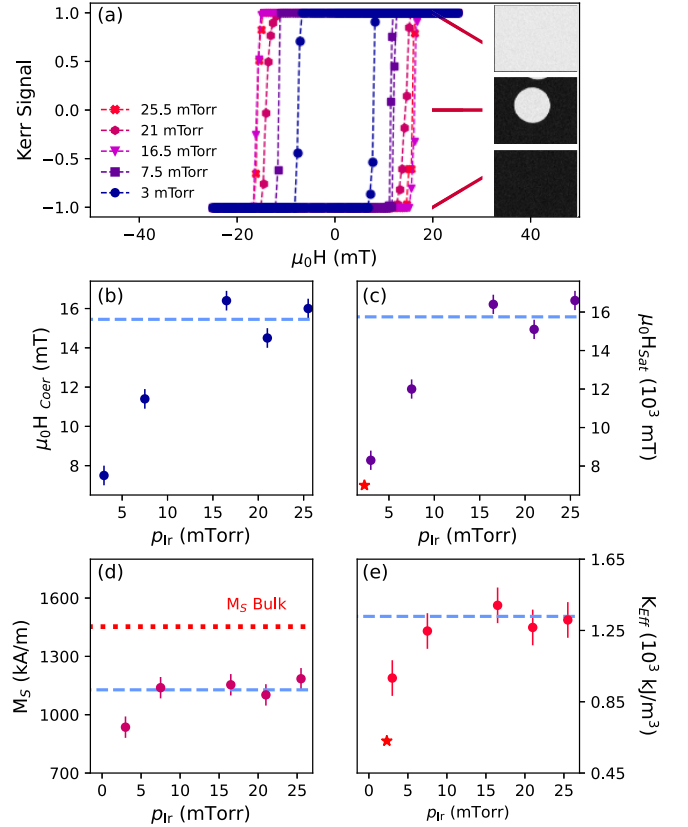


FIG. 2. Magnetometry data for the Pt/Co/Ir structures acquired at room temperature for samples where the Ir layer was grown with different Ar pressures (p_{Ir}). (a) Magnetic hysteresis cycle. The insets show typical snapshots at the indicated fields. (b) Coercive field $\mu_0 H_{\text{Coer}}$. (c) Saturation field $\mu_0 H_{\text{Sat}}$. (d) Saturation magnetization M_s . (e) Effective anisotropy constant K_{eff} . Star symbols were taken from Ref. [17]. The red dashed line in panel (d) indicate the expected bulk value for M_s . The blue dashed lines indicate the values at which the different magnitudes saturate. Results shown in panels (a), (b), and (c) were measured by PMOKE microscopy, while the results shown in panels (d) and (e) were measured using SQUID.

the XRR data does not distinguish between structural roughness and chemical intermixing between layers [13]. Although a slight decrease in χ with increasing p_{Ir} is apparent in Fig. 1(d), there is no substantial variation in the roughness-intermixing as a function of p_{Ir} . In concert with the rocking curve analysis presented in Fig. 1(c) (which demonstrated a clear increase in global roughness with increasing p_{Ir}), the unchanging nature of χ with p_{Ir} might indicate that roughness is increased and intermixing is reduced when using higher p_{Ir} values [12]. However, further characterizations, such as cross-sectional transmission electron microscopy, are required to more precisely understand the nature of the interfaces in the Pt/Co/Ir structures.

B. Static magnetic characterization

We start by describing the p_{Ir} dependence of the static magnetic properties as determined using PMOKE and SQUID via magnetization loops at room temperature with fields applied in-plane and out-of-plane, respectively. Figure 2(a) shows the

loop obtained by PMOKE microscopy, where the Kerr signal is computed as the average value of the intensity of the whole set of pixels that form each image, that in our case corresponds to an area of 0.187 mm^2 . Typical snapshots are shown in the three insets for different fields, showing that the magnetization reversal is practically dominated by the nucleation and growth of domains. From these results, the coercive and saturation fields (H_{Coer} and H_{Sat}), shown in Figs. 2(b) and 2(c), were obtained. Since the sample possesses strong perpendicular magnetic anisotropy (PMA), the values of H_{Coer} and H_{Sat} differ by a maximum of 1 mT, indicating an almost-square hysteresis loop. We observe an increase of H_{Coer} and H_{Sat} with p_{Ir} for low p_{Ir} , which tends to saturate at $p_{\text{Ir}} \approx 10 \text{ mTorr}$.

The value of M_S for a bulk Co system is 1440 kA/m , and all our samples present lower values. This phenomenology, generally observed in nanostructured systems and particularly in thin magnetic films similar to the one we are studying in this work, is usually associated with the presence of a magnetic dead layer [46–48]. Similar trends for M_S and K_{eff} in Figs. 2(d) and 2(e) emphasize that both quantities are affected by p_{Ir} in a similar way. A similar dependence with the Ar pressure within the deposition chamber was reported by Lavrijsen *et al.* [12] for M_S , K_{eff} and also about the evolution of the shape of the hysteresis cycles in Pt/Co/Pt multilayers. Lavrijsen *et al.* associate this behavior with the sensitive nature of the perpendicular magnetic anisotropy (PMA) to the growth conditions. In a similar way Wells *et al.* [13] have shown that, also for Pt/Co/Pt thin films, an increment of the Ar pressure within the deposition chamber induce a decrease of the intermixing between species of both layers, but at the same time an increment of the interface roughness. Within the context of the structural characterizations presented in Figs. 1(c) and 1(d), the initial increase in H_{Sat} , M_S , and K_{eff} with increasing p_{Ir} followed by saturation at $p_{\text{Ir}} \approx 10 \text{ mTorr}$ observed is consistent with a chemically sharper, less intermixed Co/Ir interface at higher p_{Ir} . Conversely, the similar response of H_{Coer} to p_{Ir} is consistent with an increase in the structural roughness of the Co/Ir interface with p_{Ir} [12,13,49,50].

C. Domain wall velocity under out-of-plane DC magnetic field

We now analyze the DW average velocity at room temperature under DC magnetic fields. Within this section, all the magnetic fields, to nucleate and to grow the domains, were applied in the direction perpendicular to the plane of the sample, as described in protocol (i) of Sec. III.

In Fig. 3(a) we plot $V(H)$ vs H for different p_{Ir} , as obtained from DW displacements such as those shown in the inset of Fig. 3(b). In Fig. 3(b), where we plot $\ln V(H)$ vs $(\mu_0 H)^{-1/4}$, we observe that the creep law of Eq. (1) is fairly satisfied with $\mu = 1/4$, as can be appreciated by the almost-linear relationship between the two quantities. These fits allow us to obtain the p_{Ir} dependencies of α and β , as shown in Figs. 3(c) and 3(d), respectively. For reference we also added the values of α and β obtained in a similar stack [17] for a single value of p_{Ir} .

From Figs. 3(a) and 3(b) we observe that for a given field, $V(H)$ decreases with increasing p_{Ir} , remaining constant for $p_{\text{Ir}} \geq 16.5 \text{ mTorr}$. This dependence is mainly controlled by

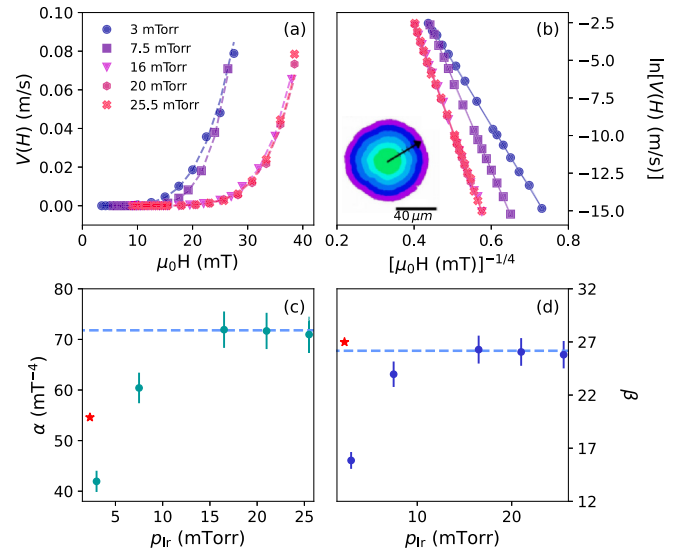


FIG. 3. (a) Velocity field characteristics $V(H)$ for different Ir growth pressures p_{Ir} . (b) Experimental measurements to the expression $\ln V(H) = -\alpha(\mu_0 H)^{-1/4} + \beta$, as expected from the creep law (1). Solid lines are linear fits to Eq. (1). Inset shows an example of the growth of a domain obtained from the sample for $p_{\text{Ir}} = 3 \text{ mTorr}$, where several pulses of $\tau = 50 \text{ ms}$ and $\mu_0 H = 10 \text{ mT}$ are applied. A typical sequence of images of a domain growing under the action of the magnetic field pulses is shown in the inset. The images correspond to the sample with $p_{\text{Ir}} = 3 \text{ mTorr}$, under magnetic field pulses of $\tau = 50 \text{ ms}$ and $\mu_0 H = 10 \text{ mT}$. From the linear fits we obtain α and β vs p_{Ir} , as shown in panels (b) and (c), respectively.

the increase of α with p_{Ir} , specially at low p_{Ir} , as shown in Fig. 3(c). Interestingly, not only α increases with p_{Ir} but also β , becoming more important at higher fields. We also observe that, for $p_{\text{Ir}} \approx 15 \text{ mTorr}$, both quantities tend to saturate. The data taken from Ref. [17] are consistent with our results and specifically with the observed trend. These p_{Ir} dependencies evince that interfacial DC magnetic characteristics and DW dynamics are correlated.

D. Asymmetric domain growth

The presence of DMI favors the formation of chiral DWs and produces an asymmetric domain growth in the presence of a perpendicular magnetic field (H) whenever a magnetic field in the direction parallel to the sample plane (H_x) is simultaneously applied. We use this method to estimate the intensity of DMI, applying the protocol (ii) of Sec. III.

Figure 4(a) shows snapshots of the domains growing in the presence of a pulsed out-of-plane field in a constant in-plane field. We observe that for $H_x = 0$ the domain growth is symmetric while a clear left-right asymmetric growth is observed for $H_x \neq 0$, as indicated. In Fig. 4(b) we show the resulting right v_{DMI}^R and left v_{DMI}^L velocities as a function of the in-plane field H_x for a given $p_{\text{Ir}} = 3 \text{ mTorr}$. The curves display clear minima at $\mu_0 H_{\text{DMI}} \sim \mu_0 |H_x| \approx 214 \text{ mT}$. In Fig. 4(c) we show H_{DMI} vs p_{Ir} and also indicate the value obtained in Ref. [17] for a similar sample.

As can be appreciated from Fig. 4(c), $\mu_0 H_{\text{DMI}} \approx 214 \text{ mT}$ and there is no appreciable and systematic change with p_{Ir}

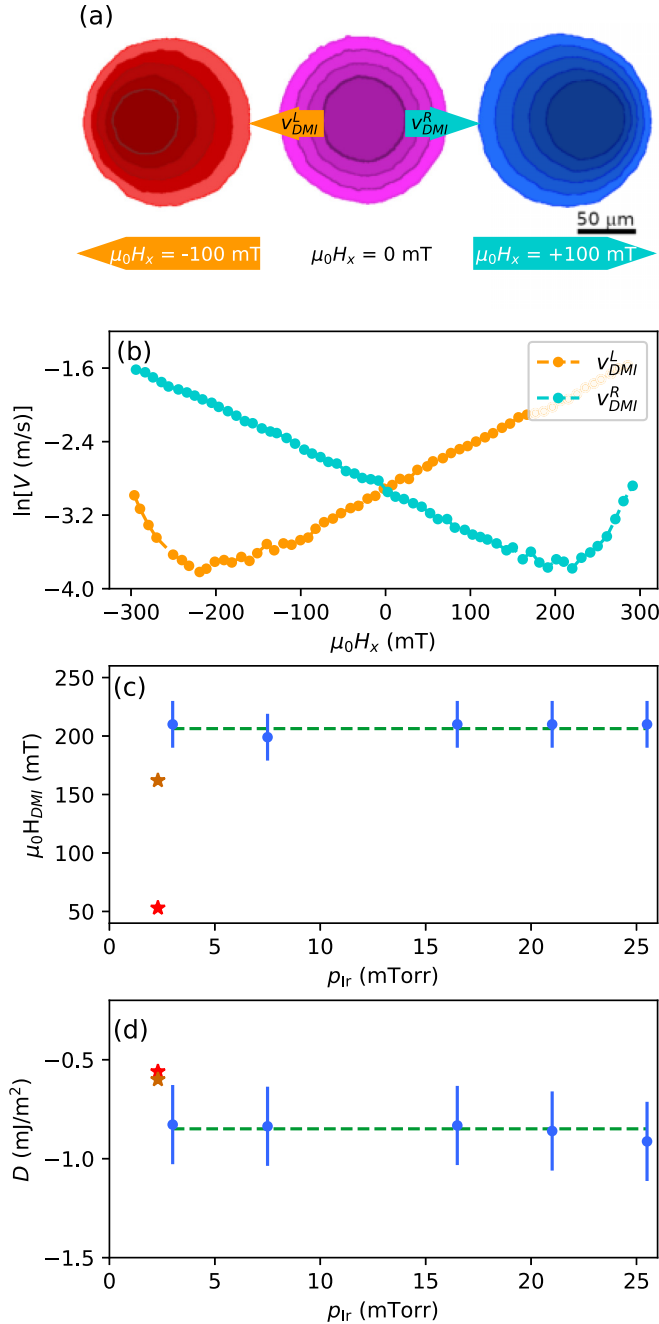


FIG. 4. (a) Domain contours grown by applying perpendicular magnetic field pulses of an amplitude of $\mu_0 H = 20$ mT with $p_{Ir} = 3$ mTorr for a constant in-plane magnetic field $\mu_0 H_x = \pm 100$ mT. (b) Typical velocities v_{DMI}^L and v_{DMI}^R measured for $\mu_0 H = 20$ mT pulses and $\mu_0 H_x = \pm 100$ mT. (c), (d) DMI field $\mu_0 H_{DMI}$ and strength $D = \mu_0 H_{DMI} M_S \sqrt{A/K_{eff}}$ as a function of p_{Ir} , using a constant $A = 17$ pJ/m. The dashed green lines indicate the average values, $\mu_0 H_{DMI} \approx 214$ mT, $D \approx -0.87$ mJ/m², and the stars are values obtained in Ref. [17], the red star from asymmetric growth and the orange from Brillouin light scattering (BLS).

in the studied range. The fact that, for the entire range of p_{Ir} of studied, we obtain an almost constant value for H_{DMI} contrasts with the trend previously discussed for M_S , K_{eff} , α , β . In Fig. 4(d) we show the corresponding DMI strength D from Eq. (3), using $A = 17$ pJ/m. We see that D is practically

constant for $p_{Ir} < 15$ mTorr, around $D \approx -0.87$ mJ/m² and for $p_{Ir} > 15$ mTorr presents a slight increase in intensity. Both in Figs. 4(b) and 4(c) we compare with results from Ref. [17] for a particular value of p_{Ir} , obtained from asymmetric growth and Brillouin light scattering.

E. AC dynamics of domain walls

The analysis of PMOKE images under alternating magnetic field pulses with zero mean is similar to the one used in the DC case (Sec. IV C) except that, in this case, after nucleation of the initial domain, pictures were taken after applying one positive square pulse followed by one negative square pulse of the same amplitude H , as described by protocol (iii) in Sec. III (from now on we refer to this pair of square pulses as the single “AC pulse”). Due mainly to the finite (negative) mean curvature of the domain and the presence of quenched disorder, the dynamics is expected to be irreversible, yielding complementary information to the DC analysis. Such information can be partially captured by monitoring simple geometrical properties such as the area, the perimeter and the roughness of the closed domain wall [19].

Figure 5(a) shows domain walls snapshots as successive AC pulses of magnetic field were applied for three samples grown at different p_{Ir} . The colors of the contours indicate the number of pulses applied before taking the picture, starting with blue contours ($N = 0$) and ending with red contours ($N = 30$). The domain evolution as a function of pulse number N is characterized by an increase of the domain-wall roughness concomitantly with the decrease of the domain area and perimeter, expected from curvature effects.

The roughening of the domain wall is shown in Fig. 5(b). We clearly observe two regimes for the mean square displacement $\langle u^2 \rangle$ as a function of the AC pulse number N , as obtained from the domain-wall contours (see Sec. II). For the five samples in Fig. 5(b) we observe that $\langle u^2 \rangle \sim N^\gamma$ with $\gamma = (0.63 \pm 0.01)$, consistent with a typical dynamic roughening process [39]. In Figs. 5(c) and 5(d) we show the evolution of the perimeter \mathcal{P} and the area \mathcal{A} , respectively, as a function of N . We observe a decrease in \mathcal{A} , as seen in previous works [18,19]. The decrease in \mathcal{P} , on the other hand, is determined by the competition between the increasing roughness, which tends to increase it and the decreasing area which tends to reduce it. We observe that both \mathcal{A} and \mathcal{P} decrease faster at larger p_{Ir} and tend to saturate on the same decay for lower pressures. Using the data of Figs. 5(c) and 5(d) and the DC velocity characteristics of Fig. 3, we can now compute the function $\Lambda(N)$ using Eq. (4). In Fig. 6(a) we show $\Delta\Lambda(N) \equiv \Lambda(N) - \Lambda(0)$ for all samples with different p_{Ir} . For each one we can fit a linear function whose slopes C are a function of p_{Ir} . From the plot of $\ln(\Delta\Lambda C^{-1})$ vs $\ln(N)$ shown in the main panel, we can see that the approximation of Eq. (4) [$\Delta\Lambda(N) \approx CN$] is valid for all p_{Ir} up to $N \approx 10$. In Fig. 6(b) we show that the fitted values of C grow with p_{Ir} with a similar behavior as H_{Coer} , H_{Sat} , M_S , and K_{eff} .

V. DISCUSSION

Our results show that the micromagnetic parameters M_S , K_{eff} , H_{Coer} , C share a similar dependence with p_{Ir} . We can see

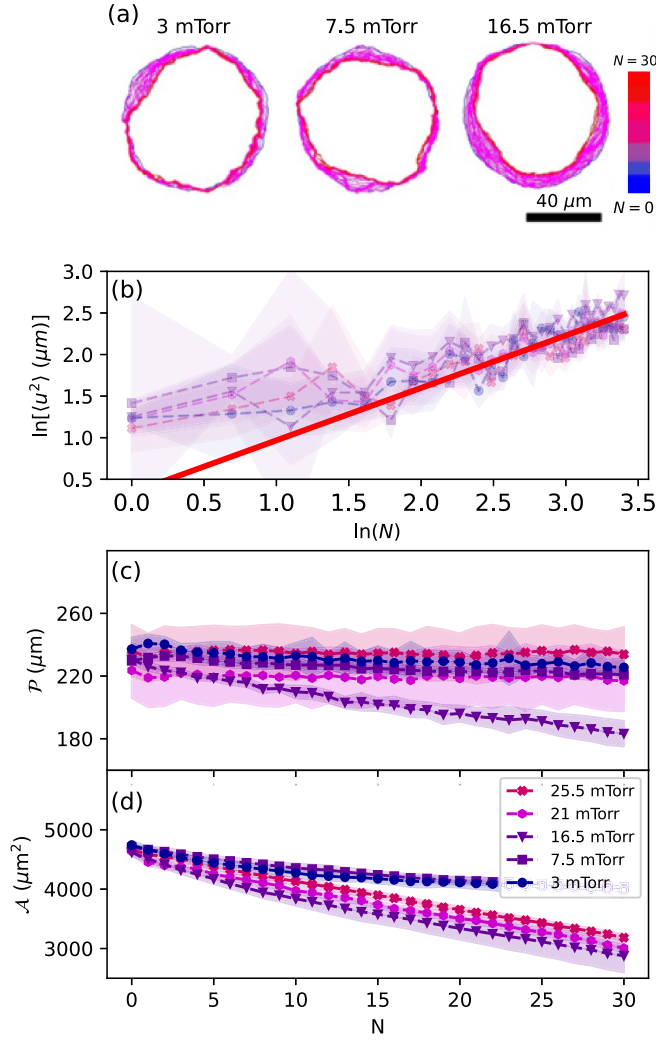


FIG. 5. (a) Contours images showing the DW shrinking and dynamical roughening as a function of the AC pulse number N , from $N = 0$ (blue contour) to $N = 30$ (red contour). Samples are grown at the different p_{Ir} as indicated. (b) Mean square displacement $\langle u^2 \rangle$ vs N for four samples grown at different p_{Ir} . The red line describes $\langle u^2 \rangle \sim N^\gamma$, with $\gamma = (0.63 \pm 0.01)$. (c), (d) Evolution of the perimeter \mathcal{P} and area \mathcal{A} of the domain as a function of N for the four samples.

that all quantities increase with increasing p_{Ir} and saturate at $p_{\text{Ir}} \approx 15$ mTorr. In all cases, the saturation value is roughly a factor 1.5 or 2 larger than the value of the same quantity at the lowest pressure $p_{\text{Ir}} = 3$ mTorr.

A similar dependence but as a function of the growth pressure of the Pt top layer was observed in Pt/Co/Pt [12]. On that case the authors attributed the observed trend to a possible crossover between two different regimes. On the one hand, a regime of considerable alloying or intermixing between Pt and Co at low pressures and, on the other hand, a regime of layered growth of the Pt layer above a given crossover pressure.

A. DC creep dynamics

Regarding the DC dynamics properties it is interesting to analyze first the dependence with p_{Ir} of the creep parameters

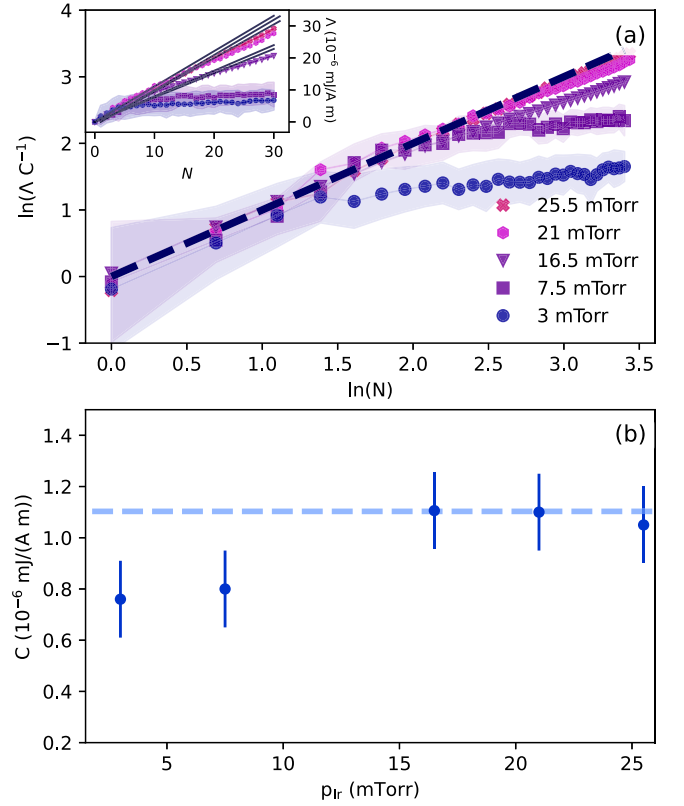


FIG. 6. (a) Rescaled function $\Delta\Lambda = \Lambda(N) - \Lambda(0)$ with C a p_{Ir} -dependent quantity. (b) Values of C obtained from the adjustments of Λ as a function of p_{Ir} , where the dashed-line is the average value of C for all samples.

α and β , defined in Eq. (1). The values of these parameters were fitted to the results obtained for the velocity as a function of the field in the creep regime. As shown in Figs. 3(b) and 3(c), these two parameters display a very similar dependence on p_{Ir} , with a saturation at $p_{\text{Ir}} \approx 15$ mTorr reaching a value which is roughly a factor 1.5 or 2 larger than the value of the same quantity at the lowest pressure $p_{\text{Ir}} = 3$ mTorr. Note that a very similar trend is observed for the micromagnetic quantities shown in Figs. 2(b)–2(e).

This behavior can be rationalized by combining the extended creep law [30] and weak-pinning theory. The extended creep law is the empirical extrapolation of the low field velocity, $v \propto e^{-(T_d/T)(H_d/H)^{1/4}}$ to larger fields, even up to H_d . The extrapolation reads $v \propto e^{-U(H)/k_B T}$ with the extended universal creep barrier $U \approx k_B T_d [(H_d/H)^{1/4} - 1]$, k_B is the Boltzmann constant, T_d is the characteristic temperature scale, and H_d is the depinning field. The extended barrier $U(H)$ hence accounts for both the divergence in the $H \ll H_d$ limit and its expected vanishing by approaching H_d . This extended though empiric form was shown to hold for a large family of thin-film materials up to H_d [30]. We can thus write

$$v \approx v_d e^{-(T_d/T)[(H_d/H)^\mu - 1]}, \quad (7)$$

from where we can identify, using Eq. (1),

$$\alpha = H_d^\mu T_d / T, \quad (8)$$

$$\beta = (T_d/T) + \ln(v_d), \quad (9)$$

In Eq. (9) we are fixing the units of velocity to m/s. Using weak collective pinning theory [20] we get

$$H_d \approx \frac{\sigma_0 \xi}{2M_S L_c^2}, \quad (10)$$

$$T_d \approx \frac{\sigma_0 \xi^2 t}{L_c}, \quad (11)$$

where L_c is the Larkin length [29,33],

$$L_c \approx \left(\frac{\sigma_0^2 \xi^2}{t^2 \Delta_0} \right)^{1/3}, \quad (12)$$

where t is the average magnetic film thickness, $\sigma_0 \approx 4\sqrt{AK_{\text{eff}}}$ is the DW energy per unit area, Δ_0 is the strength of the pinning forces, and $\xi = \max[\delta, \xi_0]$ is the correlation length of the pinning force acting on the DW. Here, ξ_0 is a constant determined by the random spatial heterogeneities. Therefore,

$$H_d \approx \left(\frac{\Delta_0^2 t^4}{8M_S \sigma_0 \xi} \right)^{1/3}, \quad (13)$$

$$T_d \approx (\Delta_0 \sigma_0 t^5 \xi^4)^{1/3}, \quad (14)$$

implying that both α and β are proportional to

$$\alpha, \beta \propto \Delta_0^{1/3} t^{5/3} (\sigma_0 \xi^4)^{1/3}. \quad (15)$$

To estimate α , we have neglected the weak contribution of H_d in comparison to T_d , as all the micromagnetic parameters of H_d in Eq. (13) are raised to the small power 1/12 or smaller. We have also neglected the contribution $\ln(v_d)$ in β because for most known samples displaying a universal creep regime we have $T_d \gg T$ and from thermal rounding experiments we know that $v_d \equiv v(H = H_d, T)$ is typically of the order of 1–10 m/s [30]. Under these assumptions a similar dependence with p_{Ir} is expected for both α and β . This estimate agrees with the results of Figs. 3(b) and 3(c).

From Eqs. (8) and (9) we can readily write

$$\beta = H_d^{-1/4} \alpha + \ln(v_d), \quad (16)$$

implying a temperature-independent linear correlation between α and β , as long as $H_d^{-1/4}$ and $\ln v_d$ do not vary appreciably for the range of α and β analyzed. In Fig. 7(a) we show that a similar linear correlation holds for α and β obtained at different p_{Ir} in our samples and also for those pairs of values α and β obtained by Shahbazi *et al.* [17] but for different Ir thicknesses. We observe that although the films have the same composition and the slopes (related to $H_d^{-1/4} T_d$) are quite similar, the intercept (related to v_d) differ appreciably. Equation (16) explains all these observations and provides an approximated method to obtain H_d and v_d . In particular, the linear fit allows us to obtain $H_d = 74 \pm 6$ mT and $v_d = 10 \pm 1$ m/s for our whole set of samples with different p_{Ir} . These values are in very good agreement with the estimate of H_d and v_d from velocity curves measured from low to high fields $H > H_d$, displaying deviations from the creep law [17]. Interestingly, such linear correlation has been also observed in several different ultrathin ferromagnetic stacks [27]. In the Appendix we show that the predicted linear correlation of Eq. (16) allows us to obtain consistent bounds for H_d and v_d for these cases. In summary, these correlations between α and

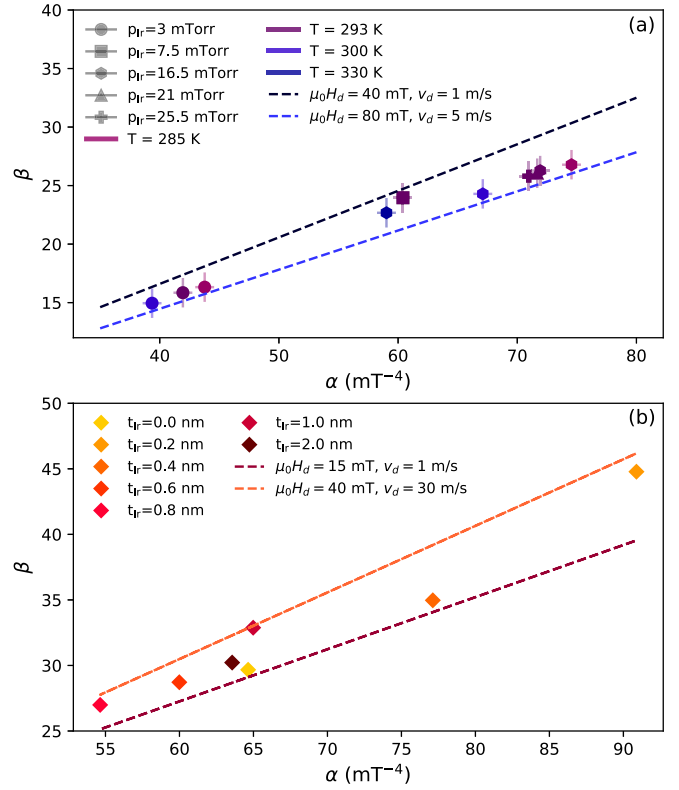


FIG. 7. Creep-law parameters β vs α for different Ta/Pt/Co/Ir/Ta samples. (a) Samples with different p_{Ir} (plot style) and temperatures (color). (b) Samples of Ref. [17] with different Ir thickness (t_{Ir}). Dotted lines indicate the bounds expected from Eq. (16) using the known range of H_d and v_d .

β are naturally explained in terms of the universal extended creep law [30].

To obtain a more microscopic insight, we analyze Eq. (15). Assuming that disorder (through Δ_0) does not depend on p_{Ir} , two situations may be considered:

(i) On the one hand, if $\xi = \xi_0 > \delta$ and that the exchange parameter A is independent of p_{Ir} , we get $\alpha, \beta \propto K_{\text{eff}}^{1/6}$, which is consistent with the results of Fig. 2(e) for K_{eff} and Figs. 3(b) and 3(c) for α and β . That is, α, β , and K_{eff} increase with increasing p_{Ir} and saturate at $p_{\text{Ir}} \approx 15$ mTorr.

(ii) On the other hand, if there is point pinning, then $\delta > \xi_0 \Rightarrow \xi = \delta \approx \sqrt{A/K_{\text{eff}}}$. In this case we get $\alpha, \beta \propto K_{\text{eff}}^{-1/2}$ which is apparently inconsistent with the p_{Ir} dependencies of the three quantities. To reconcile in this case the p_{Ir} dependencies of Fig. 2(e) for K_{eff} and Figs. 3(b) and 3(c) we need to assume that $A^{3/2}$ grows with p_{Ir} faster than $K_{\text{eff}}^{1/2}$, so that their product AK_{eff} grows with p_{Ir} .

Deciding between scenarios (i) and (ii) strictly requires a study of A vs p_{Ir} . Nonetheless, it is worth noting that scenario (i) appears to be consistent because ξ can be also estimated from Larkin weak collective pinning theory [20,38] and satisfies $\xi > \delta$. The corresponding expression is

$$\xi \approx \left[\frac{(k_B T_d)^2}{2\mu_0 H_d M_S \sigma_0 t^2} \right]^{1/3}, \quad (17)$$

TABLE I. Parameters of the full set of samples obtained through the static and dynamic characterization described in the main text.

p_{Ir} (mTorr)	3	7.5	16.5	21	25.5
M_S (kA/m)	940 (60)	1140 (60)	1150 (60)	1100 (60)	1190 (60)
$\mu_0 H_{\text{eff}}$ (T)	2.1 (0.1)	2.2 (0.1)	2.4 (0.1)	2.3 (0.1)	2.2 (0.1)
K_{eff} (kJ/m ³)	1000 (100)	1200 (100)	1400 (100)	1300 (100)	1300 (100)
δ (nm)	4.0 (0.5)	3.6 (0.5)	3.4 (0.5)	3.6 (0.5)	3.5 (0.5)
ξ (nm)	14 (7)	13 (6)	12 (6)	13 (6)	12 (6)
α ([mT] ⁻⁴)	42 (1)	60 (1)	72 (1)	72 (1)	71 (1)
β	15.8 (0.1)	24.0 (0.1)	26.3 (0.1)	26.0 (0.1)	25.8 (0.1)
$\mu_0 H_{\text{DMI}}$ (mT)	210 (20)	200 (20)	210 (20)	210 (20)	210 (20)
σ_0 (10 ⁻³ J/m ²)	16 (6)	18 (7)	19 (8)	19 (8)	19 (8)
C (10 ⁻⁶ mJ/A m)	0.5 (0.1)	0.8 (0.1)	1.1 (0.1)	1.1 (0.1)	1.0 (0.1)
σ_C (10 ⁻³ J/m ²)	0.9 (0.2)	1.8 (0.3)	2.5 (0.4)	2.4 (0.4)	2.4 (0.4)
D (10 ⁻³ J/m ²)	-0.8(0.2)	-0.8(0.2)	-0.8(0.2)	-0.9(0.2)	-0.9(0.2)

and all data for the estimate is available: From Figure 7(a) we get that $40 \text{ mT} < \mu_0 H_d < 80 \text{ mT}$. From β in Table I and Eq. (9) [neglecting again the small contribution of $\ln(v_d)$], we get $4000 \text{ K} < T_d < 7500 \text{ K}$ for the whole p_{Ir} range. From Table I we get M_S and $\sigma_0 = 4\sqrt{AK_{\text{eff}}}$ for the whole p_{Ir} range. By fixing values $A = 17 \text{ pJ/m}$ and $t \approx 0.8 \text{ nm}$ we get values of ξ which can be directly compared with the DW-width estimate $\delta \approx \sqrt{A/K_{\text{eff}}}$ vs p_{Ir} , as shown in Table I. As can be appreciated, $\xi \gtrsim 3\delta$ for the whole p_{Ir} range. One should nevertheless be cautious because this result combines several estimates, both from pinning theory and from the micromagnetic theory.

B. AC creep dynamics

We now show how the AC field creep response can yield supplementary information to the DC field creep response. In Fig. 6 we show that the effective area can also provide insights into the effects of changing p_{Ir} . The evolution of the geometrical observable Λ , defined in Eq. (4) displays two regimes as a function of the number of AC cycles N , as can be appreciated in Fig. 6(a). For small N , $\Lambda(N) - \Lambda(0) \propto CN$ as predicted by Eq. (4), allowing us to estimate the parameter C , shown in Fig. 6(b) as a function of pressure p_{Ir} . Interestingly, it displays the same trend than M_S , K_{eff} , H_{Coer} , α , β , i.e., increasing from low p_{Ir} and saturating at $p_{\text{Ir}} \approx 15 \text{ mTorr}$ (the values of all these parameters are available in Table I). As explained in Ref. [19], $C \sim \sigma_0/2M_S$ provided certain conditions are met. In particular, a homogeneous creep law response of the DW at its different points is assumed to derive Eq. (4), and this approximation was shown to be better satisfied for low values of N , both in simulations and experiments [19]. In Fig. 8 we show that the estimated elastic tension, $\sigma_C = 2M_S C$, using the measured values of C [Fig. 6(b)] and M_S [Fig. 2(d)], changes with p_{Ir} in a very similar fashion as $\sigma_0 = 4\sqrt{AK_{\text{eff}}}$ does, using $A = 17 \text{ pJ/m}$ [17] and the measured K_{eff} of Fig. 2(e). Nevertheless, we also find a whole factor ≈ 7 of difference which remains to be fully explained. One possibility is that DMI is correcting σ , as suggested by Je *et al.* [37]. The predicted correction for our case is however insufficient because it is of order $\approx 10\%$. Another possibility can be related to thickness modulations. In this case, as speculated in relation to the p_{Ir} dependencies of α and β , a possible explanation is that the effective average thickness t_{eff} seen by the dominant metastable states is systematically smaller than the magnetic

thickness t_M because they profit thinner regions to minimize the local DW energy. This would correct σ_0 effectively by $\sigma_{\text{eff}} \approx \sigma_0 t_{\text{eff}}/t_M < \sigma_0$, changing σ_{eff} in the direction of the phenomenological σ_C . Additionally, σ_C may incorporate other inhomogeneity effects which are absent in the micromagnetic bulk σ_0 . Besides the discussed difference, however, the predicted p_{Ir} dependence from the AC theory seems to be consistent with the changes of the micromagnetic parameters with p_{Ir} . The effect of p_{Ir} on M_S , H_{Coer} , K_{eff} , σ is in sharp contrast with its effect on H_{DMI} .

The AC field protocol can also yield information about the pinning potential for larger number of AC pulses N , through the analysis of the AC field induced roughening of the domain walls characterized in Figs. 5(a) and 5(b). We find that for $N \gtrsim 7$ the average quadratic width of the wall grows as $\langle u^2 \rangle \approx BN^\gamma$, with $\gamma \approx 0.63 \pm 0.01$ and B a constant prefactor. Consistently, $N \approx 7$ matches the onset of deviations from the predicted scaling $\Lambda(N) - \Lambda(0) \approx CN$ [see Eq. (4)], as shown in Fig. 6(a). This is somehow expected as such linear behavior is predicted by assuming that the domain-wall local response is homogeneous, an assumption that it is expected to hold only for the first few pulses, when the nucleated DW is still largely uncorrelated with the underlying

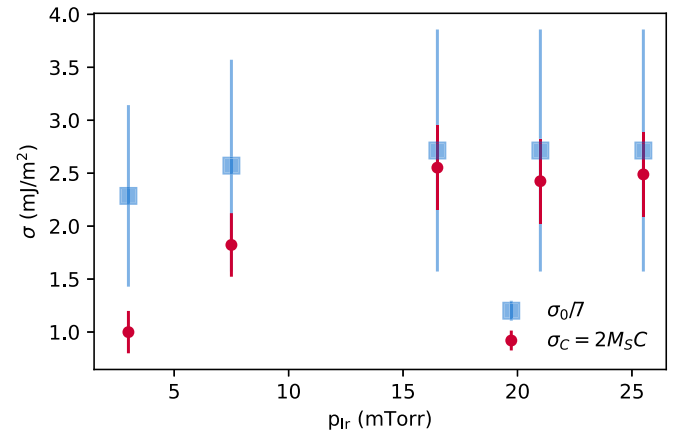


FIG. 8. Phenomenological domain wall tension σ_C obtained from the effective area $\Lambda(N)$ obtained in AC measurements, compared with the ratio $(\sigma_0/7)$ vs p_{Ir} , with $\sigma_0 = 4\sqrt{AK_{\text{eff}}}$ the micromagnetic estimate, as a function of p_{Ir} .

pinning potential. Roughening indicates, consistently, that the response becomes heterogeneous and that the domain-wall configuration gets more correlated with the underlying pinning potential as N grows. These correlations add relevant information through the exponent γ and the prefactor B . If we assume that the large-scale geometry of a DW in the AC creep regime is described by the Edwards-Wilkinson (EW) equation with an effective temperature, as it occurs in the DC creep case [24,32,51], we expect a roughening process characterized by the dynamical exponent $z = 2$ and a roughening exponent $\zeta = 1/2$, so $w_t^2 \sim t^{1/2} \sim N^{1/2}$ [39], with the effective time dictated by the pulse number N and the pulse duration τ as $t \sim N\tau$. This prediction clearly underestimates the result of 5(b) for γ . A possible explanation for the difference can be attributed to the fact that under an AC field the effective time-dependent noise acting on the domain wall is more correlated in time than the effective noise on the DC case because in the former the DW can revisit the same disorder in its oscillatory motion. This makes the effective large-scale noise time correlated or colored. If we model such colored noise $\eta(x, t)$ with an exponent $\psi > 0$, such that for two points in a DW segment $\langle \eta(x, t)\eta(x, t') \rangle \sim \delta(x - x')|t - t'|^{2\psi-1}$ ($\psi = 0$ for uncorrelated noise) linear theory predicts $z = 2$ and $\zeta = 1/2 + 2\psi > 1/2$ [39]. Therefore w_t^2 or w_N^2 should grow *faster* than in the DC case, with $\gamma = 1/2 + 2\psi > 1/2$. Our data agree with this prediction with $\psi = 0.065$ indicating a weak nonzero time correlation of the effective noise. The prefactor $B \approx \langle u^2 \rangle / N^\gamma$, on the other hand, is expected to quantify the strength of the disorder or effective noise. As can be appreciated in Fig. 5(b) the change is hardly detected for different growth pressures, thus indicating that growth pressure p_{Ir} does not have an important impact in the sample disorder. It would be interesting to test this roughening phenomenon with numerical DW simulations under AC field in random media.

VI. CONCLUSIONS

Summarizing, we have characterized the static and DC or AC dynamic properties of DWs in perpendicularly magnetized Ta/Pt/Co/Ir ultrathin films as a function of the Ir layer growth pressure. We find that the saturation magnetization, the effective anisotropy constant and the domain-wall surface tension increase monotonically with the growing pressure of the Ir layer and saturate at a threshold pressure. In sharp contrast the Dzyaloshinskii-Moriya field and the strength of the disorder remain almost constant across the same pressure range.

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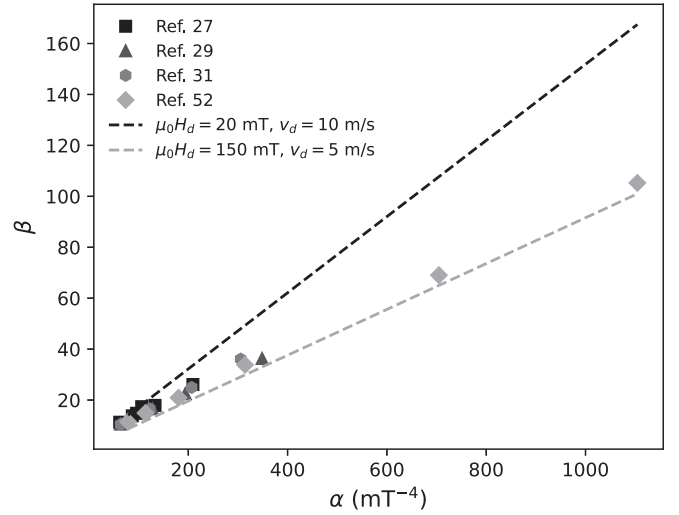


FIG. 9. Interdependence of the creep law parameters α and β obtained by Quinteros *et al.* [27], using experimental data from Refs. [29,31,52]. Using Eq. (16) we obtain values of $\mu_0 H_d$ between 20 and 150 mT, and v_d between 5 and 10 m/s.

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APPENDIX A: CREEP ANALYSIS

In Ref. [27], Quinteros *et al.* analyzed the relationship between α and β , concluding that there is a linear correlation between them, as can be appreciated in Fig. 9. We can interpret this relationship from Eq. (16), which implies that the slope $d\beta/d\alpha$ depends on H_d and the intercept on v_d .

From Refs. [29,31,52] we obtain values of $\mu_0 H_d$ between 20 and 150 mT and v_d between 10 and 5 m/s. As we can see in Fig. 9, the points are bounded by Eq. (16) with values, corresponding to ($H_d = 20$ mT, $v_d = 10$ m/s) and ($H_d = 150$ mT, $v_d = 5$ m/s). Note that the linear relation observed for high α and β values is different from the one observed at small values. This fact highlights a linear but nonuniversal relation between both quantities. In this way, instead of defining a single value for all the samples, we define a ranges of H_d and v_d that correspond to the measurements of these parameters reported in all the references. Equation (16) hence explains the apparent correlation between α and β across different samples.

APPENDIX B: SAMPLE PARAMETERS

In Table I we show the values of the different constants obtained by the static and dynamic characterizations. M_S , H_{eff} , were obtained from the hysteresis loops, K_{eff} was calculated as $\mu_0 H_{\text{eff}} M_S^2 / 2$, α and β were obtained from the lineal fits from the creep law of $V(H)$ vs $H^{-1/4}$, $\mu_0 H_{\text{DMI}}$ was obtained from the asymmetric domain growth, $\sigma_0 = 4\sqrt{AK_{\text{eff}}}$ with $A = 17$ pJ/m, C was obtained from effective area of AC dynamic, and $\sigma_C = 2 M_S C$. It is observed that the different properties of the samples increase with p_{Ir} , taking a saturation

value for $p_{\text{Ir}} > 8$ mTorr. These values are similar to those reported in some works in the bibliography [17,53]. Further-

more, although they are not the same samples, the trends correspond to what was observed by Lavrijsen *et al.* [12].

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