

Abstract

Antiferromagnets (AFs) possess unique intrinsic properties, including ultrafast dynamics, resilience to external magnetic fields, and the absence of stray-field interactions, that make them promising building blocks for next-generation spintronic and neuromorphic technologies. Yet, the controlled electrical manipulation of antiferromagnetic order and its dynamics, together with their real-space visualization, remains largely unexplored. In this thesis, synthetic antiferromagnets (SAFs) serve as model systems for real AFs, retaining essential antiferromagnetic characteristics while offering tunable interlayer coupling and experimentally accessible resonant frequencies.

The first part of this thesis demonstrates a spin-orbit torque (SOT)-driven nonvolatile memristor based on a parity-time ($\mathcal{PT}$)-symmetric SAF with perpendicular magnetic anisotropy. SOT-controlled motion of Néel-order domain walls enables reconfigurable analog memory states, while nonlinear magnetotransport provides reliable electrical readout. These results establish SAFs as promising synaptic elements for neuromorphic spintronics.

The second part explores current-driven dynamics in easy-plane SAFs via stroboscopic time-resolved magneto-optical Kerr effect (tr-MOKE) microscopy, visualizing antiferromagnetic domains in a polar MOKE geometry for the first time. Phase-resolved measurements confirm dynamic imaging of antiferromagnetic domain structures, while a strongly localized mode at a micron-scale constriction highlights geometric confinement of AF dynamics.

Finally, a chirality-controlled SOT oscillator in an easy-plane SAF is realized. Spin-rectification measurements reveal self-sustained oscillations whose handedness is set by DC current polarity, along with injection-locking to external microwave signals. Near the spin-flop transition, signatures of multiple oscillation modes and chaotic dynamics emerge.

Together, these findings establish synthetic antiferromagnet as a versatile platform for electrically controlled antiferromagnetic dynamics, spanning nonvolatile memory, gigahertz resonances, and chiral self-oscillatory states.

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Keywords: Antiferromagnetic spintronics · Synthetic antiferromagnets · Spin-orbit torques · Antiferromagnetic dynamics · Chiral self-oscillations · Nonlinear magnetotransport · Spin rectification effects · Time-resolved MOKE (tr-MOKE) microscopy

Spin-orbit driven dynamics and control of Néel order in synthetic antiferromagnets



DISSERTATION ZUR ERLANGUNG DES
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(DR. RER. NAT.)
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Das Promotionsgesuch wurde eingereicht am: 05.II.2025.
Die Arbeit wurde angeleitet von: Prof. Dr. Jörg Wunderlich.

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This research is part of the project "Cold Opto-Magnetism for Random Access Devices (COMRAD)" from the European Union's Horizon 2020 research and innovation programme under the Marie Skłodowska-Curie grant agreement No-861300. This work is further supported in part by the Deutsche Forschungsgemeinschaft (DFG, German Research Foundation) through Project-ID 314695032 – SFB 1277 "Emergente relativistische Effekte in der Kondensierten Materie".

The work described in this thesis has been carried out at the group of Prof. Dr. Jörg Wunderlich at the Institute of Experimental and Applied Physics, University of Regensburg, Universitätsstrasse 31, 93053, Regensburg, Germany.

**Spin-orbit driven dynamics and control
of Néel order in synthetic
antiferromagnets**

Pradeep Kumar Rout

*To my beloved wife Priya,
the magnetic force in my life – mostly attractive, occasionally repulsive, but always
indispensable during this long and sometimes chaotic PhD journey.*

–

*And to my amazing family and friends,
whose unwavering love and constant support have been the solid foundation
behind every step I've taken.*

*The true laboratory is the mind,
where behind illusions we uncover
the laws of truth.*

– Sir Jagadish Chandra Bose



Acknowledgements

A PhD, as many will wisely (and sometimes wearily) tell you, is a long and often bewildering journey. In my case, it has been less a solo expedition and more a collective trek – one involving many people pulling, pushing, nudging, advising, occasionally rescuing, and frequently caffeinating me along the way. I hope I have acknowledged everyone who contributed; if I missed someone, please accept this all-purpose, oversized **Thank You** in advance. To the readers of this thesis: thank you for investing your time in exploring these pages. I hope you find something enjoyable, enlightening, or at least mildly entertaining amid the equations and graphs. For questions, comments, or critical remarks, feel free to reach out – preferably during daylight hours.

The most important acknowledgement is for my supervisor, [Prof. Dr. Jörg Wunderlich](#). Our paths crossed at a magnetism conference at NISER, Bhubaneswar, where a discussion over my poster on *"magnetotransport effects in non-collinear antiferromagnets"* led him to offer me a PhD position, an assessment for which I am still immensely grateful. Working on antiferromagnetic spintronics under the guidance of one of the field's pioneers was a dream come true. [Jörg](#) is the friendliest professor I have met, approachable, encouraging, and passionate about physics. He eagerly discusses results, brainstorms ideas, and unravels physics mysteries with genuine curiosity. While his bureaucratic skills might need a firmware update, his mentorship, kindness, and scientific insight more than compensate.

Much of my PhD became collaborative, especially after [Jörg](#) moved to Regensburg. Many experiments could not have been performed in our own labs. I am deeply grateful to the group of [Prof. Dr. Christian Back](#) at the Technical University of Munich for providing access to their time-resolved magneto-optical Kerr effect (tr-MOKE) microscopy setup, which was crucial for [chapter V](#). Special thanks to

Dr. Franz Vilmeier, and Dr. Christian Riedel for their expertise, guidance, and patience whenever I arrived in Munich.

High-quality samples were essential for our work. I sincerely thank Prof. Dr. Olav Hellwig and Dr. Ruslan Salikhov (HZDR, Dresden and TU Chemnitz) for providing synthetic antiferromagnets (SAFs) precisely tailored to our needs, and for the discussions, draft reviews, and motivation to stay curious. I am also grateful to Zbyněk Soban (FZU, Prague) for tirelessly fabricating nanoconstriction devices and Hall bars, and I admit that my requests for more samples and devices will probably continue for a few more years. Special thanks to Rubén Otxoa (Hitachi Cambridge Laboratory) for his guidance in simulating the dynamics of antiferromagnetically coupled systems.

Next, I would like to thank Prof. Dr. Jaroslav Fabian, Prof. Dr. Sascha Schäfer, and Prof. Dr. Ferdinand Kuemmeth for serving on my doctoral committee and for the time and care they dedicated to carefully reading and providing their valuable feedback on this dissertation.

I am grateful to all members of the **Wunderlich group**, (now COMPAS: Compensated Magnets - Properties & Advanced Spintronics group). João Godinho, I learned immensely from you and enjoyed every moment of working together. Special thanks to Matthias Kronseder, Magdalena Okon, Markus Hollenberger, Anna Girnghuber (now Schmid), Carina Maue, Tomáš Janda and Jungwoo Koo for their support, friendship, and for tolerating me for all these years. You will probably have to endure a bit more of my complaints about the German Mensa food; nevertheless, you were truly awesome. You created the best coworking atmosphere in the lab and an even better friendship outside it. As if thanking you once were not enough, thank you all again for reading this thesis and for your comments, which prevented it from being objectively worse. Special thanks to Anna for translating the abstract into German and giving the illusion of my German language competence. Thanks to Jay and family for the mouthwatering Korean dinner, to Anna and Carina for the delicious cakes and cookies, and I promise to provide more tandoori chicken after finishing the thesis.

My early PhD days at the Institute of Physics (FZU), Czech Academy of Sciences, in Prague, in the group of Prof. Tomas Jungwirth, were enlightening. Thanks to all colleagues there: Prof. Tomas Jungwirth, Prof. Vit Novak, Dr. Kamil Olejnik, Dr. Jakub Železný, Dr. Karel Výborný, Dr. Xavi Marti, Dr. Phillip Křížek, Dr. Zbyněk

Soban, Dr. Jan Žubák, Dr. Zdeněk Kašpar, Dr. João Godinho, Dr. Srikrishna, and especially Eva Matuszyna, my saviour and best friend who fearlessly tackled administrative chaos. I will always cherish the Pilsner Urquell, Smažený, and koláč shared with all of you.

As part of the Marie Skłodowska-Curie ITN COMRAD consortium, I traveled across Europe for workshops and meetings, forming friendships with Aya, Julian, David, Paul, Jackson, Timur, Thomas, Theodoros, Guillermo, Lukas, Nikolay and Leo, with whom physics discussions and evening outings created invaluable memories.

I thank my Master's thesis supervisor, Prof. Dr. Ajaya Kumar Nayak (NISER, Bhubaneswar), for fostering my curiosity and motivating me academically, laying the foundation for my PhD journey.

I am grateful to my family in India, my parents, in-laws, brothers, and sister-in-law, for their love and encouragement. During the writing of this thesis, I became an uncle to my newborn nephew, Kirtiman, whose arrival brought great joy and happiness as we welcomed the newest member of our family. Thanks also to my friends worldwide and my "family away from home" in Regensburg and Prague: Narendra, Prateek, Thanuj, Albert, Souvik, Shaon, Manuela, Laura, Barno, Hande, Debam, Bashab, Nithin, Suman, Srikrishna, Yuta, Koen, Joao, Tomas, Jay, Anna, Abinash, Supriti, Ashish, Stiti, Saroj, Sarita, Alok, Manas, Mahashweta, Anjan, Shilpa, Ashok, Akshatha, Animesh, Soumya, Vivek, Vivek, Binoy, Bimalesh, Charanpreet, Dola, Subir, Mona, Sanjoy, Dulal, Bipul, Sipun, Kishan, Arun, Abhaya, Supriya, Jaydev.

Finally, I dedicate a special note to my wife, Priya, my childhood friend and life partner. The most amazing thing during my PhD is that I married you. Your love, patience, and support have been my anchor, making even the toughest days brighter. I am endlessly grateful for sharing this journey with you.

Thank you all for your support, friendship, and countless memories that made this journey not only possible but truly unforgettable.

Abstract

Antiferromagnets (AFs) offer ultrafast dynamics, robustness against magnetic fields, and the absence of stray-field interactions, attributes that position them as promising building blocks for next-generation spintronic and neuromorphic technologies. Yet, the controlled electrical manipulation of antiferromagnetic order and its dynamics, along with their direct real-space visualization, remains largely unexplored. In this thesis, we employ synthetic antiferromagnets (SAFs) as model systems for real AFs, which preserve key antiferromagnetic properties while providing tunable interlayer coupling and accessible resonant frequencies for systematic studies.

The thesis presents three main experimental contributions. First, we demonstrate a spin-orbit torque (SOT)-driven nonvolatile memristor based on a parity-time ($\mathcal{PT}$)-symmetric SAF with perpendicular magnetic anisotropy (PMA). SOT controlled Néel-order domain wall (DW) motion enables reconfigurable, analog memory states, while nonlinear magnetotransport provides robust electrical readout, establishing compensated SAFs as promising synaptic elements for neuromorphic spintronics. The absence of long-range stray fields ensures stable, reproducible, and electric current pulse-driven control of domain configurations. These results underscore the potential of SAF as a versatile nonvolatile memory that mimics an artificial synapse, with applications in neuromorphic computing, where the reconfigurable synaptic weight is encoded in the ratio of reversed antiferromagnetic domains.

Second, in in-plane easy-plane SAFs, we use stroboscopic time-resolved magneto-optical Kerr effect (tr-MOKE) microscopy to visualize, for the first time, antiferromagnetic domains in real space using polar magneto-optical Kerr effect (p-MOKE), which is otherwise inaccessible for this system in a p-MOKE geometry. The p-MOKE contrast is enabled by a staggered damping-like field arising from

current-induced SOT acting oppositely on the two magnetizations of the SAF, resulting in out-of-plane oscillations of the magnetizations. By performing phase-resolved tr-MOKE measurements, we uncover that the Néel order in adjacent domains oscillates 180° out of phase, further corroborating our dynamic visualization of antiferromagnetic domains using polar tr-MOKE. In addition, DC currents enable reversible DW motion, which is directly observed dynamically. Finally, we demonstrate the existence of a strongly localized antiferromagnetic mode confined to the micron sized constriction, highlighting the spatial confinement of antiferromagnetic dynamics imposed by the device geometry.

Finally, we realize a chirality-controlled SOT-oscillator in an easy-plane SAF. Spin rectification measurements and simulations demonstrate current-driven self-sustained oscillations, whose handedness is determined by the DC current polarity and which can injection-lock to external microwave signals. The experiments further indicate the onset of multiple oscillation modes and chaotic dynamics near the spin-flop transition. To our knowledge, these results provide the first experimental indications consistent with SOT-driven self-oscillations in an AF.

Together, these results establish SAF as a powerful platform for electrically controlled antiferromagnetic dynamics, spanning nonvolatile memory operation, gigahertz resonances, and self-oscillatory states. The demonstrated mechanisms and device concepts offer promising routes toward ultrafast, energy-efficient spintronic and neuromorphic hardware and pave the way for translating similar functionalities into crystalline AFs and altermagnets.

Keywords: antiferromagnetic spintronics, synthetic antiferromagnets (SAFs), spin-orbit torques (SOTs), antiferromagnetic dynamics, chiral self-oscillations, nonlinear magnetotransport, time-resolved magneto-optical Kerr effect (tr-MOKE) microscopy, spin rectification effects (SREs).

Zusammenfassung

Antiferromagneten (AFs) zeichnen sich durch ihre ultraschnelle Dynamik, hohe Robustheit gegenüber Magnetfeldern und das Fehlen von Streufeldwechselwirkungen aus – Eigenschaften, die sie zu vielversprechenden Bausteinen für Spintronik- und neuromorphe Technologien der nächsten Generation machen. Dennoch sind die kontrollierte elektrische Manipulation der antiferromagnetischen Ordnung und Dynamik sowie deren direkte Visualisierung im realen Raum bislang weitgehend unerforscht. In dieser Arbeit verwenden wir synthetische Antiferromagneten synthetic antiferromagnets (SAFs) als Modellsysteme, die zentrale antiferromagnetische Eigenschaften bewahren und gleichzeitig eine einstellbare Interlagenkopplung sowie zugängliche Resonanzfrequenzen für systematische Untersuchungen bieten.

Die Arbeit umfasst drei wesentliche experimentelle Beiträge: Zu Beginn demonstrieren wir einen durch Spin-Bahn-Drehmoment (englisch spin-orbit torque (SOT)) angetriebenen nichtflüchtigen Memristor auf der Grundlage eines Paritäts- und Zeitumkehr ($\mathcal{PT}$)-symmetrischen SAF mit senkrechter magnetischer Anisotropie. Die durch SOT induzierte, kontrollierte und reversible Propagation von 180° -Néel-Domänenwänden ermöglicht rekonfigurierbare, analoge Speicherzustände. Der nichtlineare Magnetotransport erlaubt ein direktes elektrisches Readout des Speicherzustands, das dem Verhältnis antiferromagnetischer Domänen mit entgegengesetztem Néel-Vektor entspricht. Damit erscheinen kompensierte SAFs als geeignete Kandidaten für synaptische Bauelemente in der neuromorphen Spintronik. Das Fehlen weitreichender Streufelder ermöglicht eine stabile, reproduzierbare und durch elektrische Strompulse gesteuerte Kontrolle der Domänenkonfigurationen. Diese Ergebnisse unterstreichen das Potenzial von SAFs als vielseitige nichtflüchtige Speicher, die künstliche Synapsen nachahmen und in der neuromorphen Datenverarbeitung An-

wendung finden, wobei das rekonfigurierbare synaptische Gewicht im Verhältnis der umgekehrten antiferromagnetischen Domänenstruktur kodiert ist.

zweiten Teil meiner Doktorarbeit charakterisieren wir kompensierte SAFs mit Easy-Plane-Magnetanisotropie unter SOT-Anregung mittels stroboskopischer, zeitaufgelöster magneto-optischer Kerr-Mikroskopie (tr-MOKE) und machen damit erstmals antiferromagnetische Domänen im Ortsraum mittels polarem MOKE sichtbar, was in dieser Geometrie sonst nicht möglich wäre. Der p-MOKE-Kontrast entsteht durch ein effektives, gestaffeltes antidämpfungsartiges Spin-Bahn-Feld, das Oszillationen der antiferromagnetischen Ordnung anregt, bei denen – abhängig von der lokalen Domänenkonfiguration – eine dynamische und rechtwinklig zur Probenebene gerichtete Netto-Magnetisierungskomponente entsteht. Phasenaufgelöste MOKE-Messungen zeigen, dass die Néel-Ordnung in benachbarten Domänen um 180° phasenverschoben oszilliert und bestätigen so die dynamische Visualisierung antiferromagnetischer Domänen mittels polarer tr-MOKE. Darüber hinaus führen Gleichströme zu einer reversiblen Domänenwandbewegung, die direkt beobachtet werden kann. Schließlich identifizieren wir eine stark lokalisierte antiferromagnetische Mode, die nur innerhalb einer Nanoconstriction generiert wird.

Abschließend zeigen wir die Realisierung eines chiralen antiferromagnetischen SOT-Oszillators in einem Easy-Plane-SAF, dessen Chiralität durch die Polarität des den Oszillator antreibenden Gleichstroms bestimmt wird. Spin-Gleichrichtungsmessungen und Simulationen belegen Autooszillationen, die sich bei zusätzlich angelegten monochromatischen Mikrowellensignalen in Form von Injection Locking an diese anpassen. Zudem beobachten wir das Auftreten mehrerer Schwingungsmodi sowie chaotische Dynamik nahe des Spin-Flop-Übergangs. Nach unserem Verständnis handelt es sich dabei um die erste experimentelle Beobachtung SOT-getriebener Autooszillationen in einem Antiferromagneten.

Zusammengefasst zeigen diese Ergebnisse, dass synthetische Antiferromagneten als leistungsstarke Plattform für elektrisch gesteuerte antiferromagnetische Dynamiken, nichtflüchtige Speicheroperationen, Gigahertz-Resonanzen und selbstoszillierende Zustände dienen können. Die demonstrierten Mechanismen und Gerätekonzepte bieten vielversprechende Wege zu ultraschneller, energieeffizienter spintronischer und neuromorpher Hardware und ebnen den Weg für die Übertragung ähnlicher Funktionen auf kristalline Antiferromagneten und Altermagnete.

Schlüsselwörter: Antiferromagnetische Spintronik, Synthetische Antiferromagneten, Spin-Bahn-Drehmoment, Antiferromagnetische Dynamik, Chirale Selbstschwingungen, Nichtlinearer Magnetotransport, zeitaufgelöste magneto-optische Kerr Effekt (tr-MOKE) Mikroskopie, Spin-Gleichrichtungs-Effekte.

ସାରାଂଶ

ପ୍ରତି ସମାନ୍ତରଳ ରୁମ୍ବକୀୟ ଧାତୁଗୁଡ଼ିକ (ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍) ହୁଏ ଗତିଶୀଳତା, ରୁମ୍ବକୀୟ କ୍ଷେତ୍ର ପ୍ରତି ଦୃଢ଼ତା, ଏବଂ ବିପଥଗାମୀ ରୁମ୍ବକୀୟ କ୍ଷେତ୍ରର ଅନୁପସ୍ଥିତି ପ୍ରଦାନ କରିଥାନ୍ତି । ଏହି ଗୁଣାବଳୀ ସେମାନଙ୍କୁ ପରବର୍ତ୍ତୀ ପିଢ଼ୀର ସ୍ପିନ୍ଟ୍ରେନିକ୍ସ ଏବଂ ନ୍ୟୁରୋମର୍ଫିକ୍ ପ୍ରଯୁକ୍ତିର ପ୍ରତିଷ୍ଠାକାରୀ ମୌଳିକ ଉପାଦାନ ଭାବେ ପ୍ରତିଷ୍ଠା କରେ । ତଥାପି, ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ କ୍ରମର ନିୟନ୍ତ୍ରିତ ବିଦ୍ୟୁତିକ ସଞ୍ଚାଳନ ଏବଂ ତାହାର ଗତିବିଧି, ସହିତ ସେମାନଙ୍କର ସିଧାସଳଖ ବାସ୍ତବ ଅବକାଶ ଦୃଶ୍ୟାବଳୀ, ମୁଖ୍ୟତଃ ଅନୁସନ୍ଧାନ ହୋଇନି । ଏହି ଗବେଷଣା ପ୍ରବନ୍ଧରେ, ଆମେ କୃତ୍ରିମ ପ୍ରତି ସମାନ୍ତରଳ ରୁମ୍ବକୀୟ ଧାତୁ (SAF) କୁ ନମୁନା ତନ୍ତ୍ର ଭାବେ ବ୍ୟବହାର କରୁଛୁ, ଯାହା ମୁଖ୍ୟ ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ଗୁଣ ରକ୍ଷା କରେ ଏବଂ ସହଜ ପରତଃ ସ୍ତର ସଂଯୋଗ ଓ ଉପଲବ୍ଧ ଅନୁକମ୍ପନ ଆବୃତ୍ତି ପ୍ରଦାନ କରିଥାଏ, ଯାହା ପ୍ରଣାଳୀବଦ୍ଧ ଅଧ୍ୟୟନ ପାଇଁ ଉପଯୋଗୀ ।

ଏହି ଗବେଷଣା ପ୍ରବନ୍ଧରେ ତିନିଟି ପ୍ରମୁଖ ପ୍ରୟୋଗାତ୍ମକ ଅବଦାନ ପ୍ରଦର୍ଶିତ କରାଯାଇଛି । ପ୍ରଥମେ, ଆମେ ପ୍ରଦର୍ଶିତ କରିଛୁ ଯେ ସମତା ସମ୍ମିଳିତ ଏବଂ ସମୟ ପ୍ରତିକ୍ରିୟା ସମ୍ମିଳିତ ଏବଂ ଲମ୍ବଗତ ରୁମ୍ବକୀୟ ବିଦିଗତ ଅସମତା ଥିବା SAF ଉପରେ ଏକ ସ୍ପିନ୍ ଓରବିଟ୍ ଟର୍କ୍ (SOT) ଚାଲିତ ସ୍ଥିର ମେମ୍ବ୍ରିଷ୍ଟର୍ ନିର୍ମାଣ କରାଯାଇପାରେ । SOT ନିୟନ୍ତ୍ରିତ ନୀଲ୍ ଅର୍ଡର ଡୋମେନ୍ ଝାଲୁ (DW) ର ଗତି ପୁନଃବିନ୍ୟାସଯୋଗ୍ୟ, ଏକାଧାର୍ ସ୍ଥାପିତ (ମେମୋରୀ) ସ୍ଥିତିକୁ ସକ୍ଷମ କରିଥାଏ, ଯେଉଁଠାରେ ନିମ୍ନିଅର୍ ମ୍ୟାଗ୍ନେଟୋଟ୍ରାନ୍ସପୋର୍ଟ ଦୃଢ଼ ବୈଦ୍ୟୁତିକ ପ୍ରକାଶ ପ୍ରଦାନ କରେ — ଏହା SAF କୁ ନ୍ୟୁରୋମର୍ଫିକ୍ ସ୍ପିନ୍ଟ୍ରେନିକ୍ସ ପାଇଁ ପ୍ରତିଷ୍ଠାକାରୀ ସିନେପ୍ଟିକ୍ ଉପାଦାନ ଭାବେ ପ୍ରତିଷ୍ଠା କରେ ଦୀର୍ଘ ଦୂର ବିପଥଗାମୀ ରୁମ୍ବକୀୟ କ୍ଷେତ୍ରର ଅଭାବ ଡୋମେନ୍ ବିନ୍ୟାସର ସ୍ଥିର, ପୁନଃରୁପାଦନୀୟ, ଏବଂ ବୈଦ୍ୟୁତିକ ପଲ୍ଲ ଚାଲିତ ନିୟନ୍ତ୍ରଣ ନିଶ୍ଚିତ କରେ ଏହି ପରିଣାମଗୁଡ଼ିକ SAF କୁ ଏକ ବହୁ—ପ୍ରୟୋଗୀ ନନ୍—ଭୋଲଟାଜିକ୍ (ସ୍ଥିର) ମେମୋରୀ ଭାବେ ପ୍ରତିଷ୍ଠା କରେ, ଯାହା କୃତ୍ରିମ ସିନାପ୍ସକୁ ଅନୁକରଣ କରେ । ଏଠାରେ ପୁନଃବିନ୍ୟାସଯୋଗ୍ୟ ସିନେପ୍ଟିକ୍ ଓଜନ ବିପରୀତ ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ଡୋମେନ୍ ଅନୁପାତରେ ଏନକୋଡ୍ ହୋଇପାରିବ ।

ଦ୍ୱିତୀୟତଃ, ଇନ—ପ୍ଲେନ୍ ଇଜି—ପ୍ଲେନ୍ SAF ରେ, ଆମେ ଷ୍ଟ୍ରୋବୋସ୍କୋପିକ୍ ଟାଇମ୍ ରିଜଲୁସ୍ ମ୍ୟାଗ୍ନେଟୋ ଅପ୍ଟିକାଲ୍ କେର୍ ଇଫେକ୍ଟ୍ (tr MOKE) ମାଇକ୍ରୋସ୍କୋପି ବ୍ୟବହାର କରି, ପ୍ରଥମେ ବାସ୍ତବ ସ୍ଥାନରେ ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ଡୋମେନ୍ ଗୁଡ଼ିକର ସଙ୍ଗତିଶୀଳ ଗିଗାହର୍ଜ୍ ଦୋଳନ ଦେଖାଇଛୁ । ଏହି ଦୋଳନ ବିଦ୍ବୁତ୍ ଚାଲିତ SOT ଦ୍ୱାରା ଚାଲିତ ଏବଂ ପ୍ରକାଶ କରେ ଯେ ପାଶ୍ଚାତର୍ତ୍ତୀ ଡୋମେନ୍ ମଧ୍ୟରେ ନୀଲ୍—କ୍ରମ ୧୮୦° ଫେଜ୍ ଅନ୍ତରରେ ଦୋଳନ କରେ ତତ୍ତ୍ୱପରି, ସିଧାସଳଖ

ବୈଦ୍ୟୁତିକ କରେଣ୍ଟ ଚାଳିତ ରିଭର୍ସିବଲ୍ ଡୋମେନ୍ ଡିଭାଇସ୍ ଗଠିତ କରାଯାଇଥିବା MOKE ମାଧ୍ୟମରେ ପ୍ରତ୍ୟକ୍ଷ ଭାବେ ଗତିଶୀଳ ଭାବରେ ଦେଖାଯାଏ । ଏହି ଗବେଷଣାରେ, ଆମେ ପ୍ରଦର୍ଶନ କରିଛୁ ଯେ ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ଅନୁକମ୍ପନ ଜ୍ୟାମିତିକ ସଂକୋଚନରେ ଗୁରୁତ୍ୱପୂର୍ଣ୍ଣ ଭାବରେ ସୀମିତ, ଯାହା ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ଅନୁକମ୍ପନର ଜ୍ୟାମିତିକ ସ୍ଥାନାନ୍ତରଣକୁ ପ୍ରକାଶ କରେ ।

ଶେଷରେ, ଆମେ ଇତି—ପ୍ଲେନ୍ SAF ରେ ଏକ ଧୀରତା ନିୟନ୍ତ୍ରିତ SOT ଦ୍ୱାରା ପ୍ରଚଳିତ ଦୋଳନ ସୃଷ୍ଟି କରିଛୁ । ସ୍ପିନ୍ ରେକ୍ଟିଫିକେସନ୍ ମାପନ ଏବଂ ଅନୁକରଣ ଦ୍ୱାରା କରେଣ୍ଟ ଚାଳିତ ସ୍ୱୟଂ—ଦୋଳନ ଦେଖାଯାଇଛି, ଯାହାର ସ୍ଥାୟୀତ୍ୱ ନେତୃତ୍ୱ ସିଧାସଳଖ ବୈଦ୍ୟୁତିକ କରେଣ୍ଟ ଚିହ୍ନ ଦ୍ୱାରା ନିର୍ଣ୍ଣୟ ହୁଏ ଏବଂ ଯାହା ବାହ୍ୟ ମାଲକ୍ତ୍ରାସ୍ତ୍ର ସଙ୍କେତ ସହିତ ଇଞ୍ଜେକ୍ସନ୍—ଲକ୍ କରାଯାଇପାରେ । ପ୍ରୟୋଗଗୁଡ଼ିକ ଅନେକ ଦୋଳନ ମୋଡ୍ ଏବଂ ସ୍ପିନ୍—ଫ୍ଲପ୍ ସଂକ୍ରମଣ ନିକଟରେ ବିଶୁଦ୍ଧିତ ଗତିବିଧି ଆରମ୍ଭ ସୂଚିତ କରେ । ଆମ ଜ୍ଞାନାନୁସାରେ, ଏହା ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ରେ SOT ଚାଳିତ ସ୍ୱୟଂ—ଦୋଳନ ଦେଖାଯିବା ପ୍ରଥମ ପରିବେଷଣ ।

ସମସ୍ତ ଏହି ପରିଣାମ SAF କୁ ବିଦ୍ରୁତ ନିୟନ୍ତ୍ରିତ ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ଗତିବିଧି ପାଇଁ ଏକ ଶକ୍ତିଶାଳୀ ପ୍ଲଟଫର୍ମ ଭାବେ ସ୍ଥାପିତ କରେ, ଯାହା ସ୍ଥିର ମେମୋରୀ ଚାଳନା, ଗିଗାହର୍ଜ୍ ଅନୁକମ୍ପନ, ଏବଂ ସ୍ୱୟଂ—ଦୋଳନ ଅବସ୍ଥାକୁ ଆବର୍ତ୍ତ କରେ । ପ୍ରଦର୍ଶିତ ପ୍ରଣାଳୀ ଏବଂ ଡିଭାଇସ୍ ଧାରଣା ଦ୍ରୁତ, ଉର୍ଦ୍ଧ୍ୱାଦକ୍ଷ ସ୍ପିନ୍ଟ୍ରୋନିକ୍ ଏବଂ ନ୍ୟୁରୋମର୍ଫିକ୍ ହାର୍ଡୱେର୍ ପାଇଁ ପ୍ରତିଷ୍ଠାକାରୀ ମାର୍ଗ ଦେଖାଇଥାଏ ଏବଂ ସମାନ କାର୍ଯ୍ୟକାରୀତା ସ୍ତରକୁ ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ଏବଂ ଅଲ୍ଟରମ୍ୟାଗ୍ନେଟିକ୍ରେ ଅନୁବାଦ କରିବା ପାଇଁ ମାର୍ଗ ପ୍ରସ୍ତୁତ କରେ ।

ମୁଖ୍ୟ ଶବ୍ଦଗୁଡ଼ିକ: ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ସ୍ପିନ୍ଟ୍ରୋନିକ୍ସ, କୃତ୍ରିମ ପ୍ରତି ସମାନ୍ତରଳ ରୁଫ୍‌କାୟ ଧାତୁ (SAF), ସ୍ପିନ୍—ଓରବିଟ୍ ଟର୍କ୍ (SOT), ଆଣ୍ଟିଫେରୋମ୍ୟାଗ୍ନେଟିକ୍ ଗତିବିଧି, ଧୀରତା ନିୟନ୍ତ୍ରିତ ସ୍ୱୟଂ—ଦୋଳନ, ନିଉନିଅର୍ ମ୍ୟାଗ୍ନେଟୋଟ୍ରାନ୍ସପୋର୍ଟ, ସ୍ପିନ୍ ରେକ୍ଟିଫିକେସନ୍ ପ୍ରଭାବ ।

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Introduction and thesis trajectory

*Antiferromagnetic materials are internally magnetic, yet their alternating atomic magnetic moments cancel out, rendering them externally invisible. This unique property makes information stored in antiferromagnets (AFs) largely insensitive to magnetic fields, ideal for ultra-dense, interference-free spintronic devices. Their complex ultrafast dynamics further set them apart from ferromagnets (FMs). This chapter introduces the field of antiferromagnetic spintronics ([section 1.1](#)), outlines recent breakthroughs, and establishes the context for the experimental studies presented in this thesis *. Synthetic antiferromagnets (SAFs), the primary systems examined in this thesis, are presented as versatile and readily accessible model platforms for investigating antiferromagnetic phenomena in [section 1.2](#). The emerging direction of neuromorphic spintronics is briefly discussed in [section 1.3](#), with emphasis on the key challenges and opportunities ahead. The chapter concludes with a concise summary of the thesis' main experimental findings and an outline of the chapters that follow ([section 1.4](#)).*

Keywords: antiferromagnetic spintronics, neuromorphic spintronics, synthetic antiferromagnets (SAFs), spin-orbit torques (SOTs), and antiferromagnetic dynamics.

*This chapter motivates the thesis and situates it within the existing literature, while necessarily introducing some technical terms whose detailed discussion follows in subsequent chapters.

1.1 Antiferromagnets in spintronics: background and state of the art

“They are extremely interesting from the theoretical viewpoint, but do not seem to have any applications”.¹ This paradigm, clearly articulated in the 1970 Nobel lecture by Louis Néel, reflected the prevailing perception of AFs as scientifically intriguing yet practically inert materials. As he emphasized, their lack of net magnetization made them unsuitable for applications that relied on conventional magnetic manipulation, such as those involving FMs. However, in recent years, this perception has undergone a profound shift. AFs have emerged as active components in spintronic research, no longer merely passive systems but rather functional materials with unique advantages such as ultrafast dynamics, robustness against external magnetic fields, and scalability. This evolution has led to the formulation of *antiferromagnetic spintronics* – a field that, as some have aptly put it, *“makes antiferromagnets useful and spintronics more interesting”*^{2–6}.

Although FMs have long been known and widely utilized due to their distinct magnetic behavior, AFs, despite being more prevalent in nature, have only recently attracted significant research interest^{7,8}. This delayed attention can be attributed to the lack of characteristic features in AFs that make FMs so easily identifiable and functional, such as the presence of stray magnetic fields and their pronounced reaction to applied magnetic fields. These observable properties in FMs stem from the uniform, parallel orientation of atomic magnetic moments, which collectively give rise to a net magnetization. As depicted in [figure 1.1a](#), this alignment within a cubic unit cell results in a coherent and macroscopically detectable magnetic field, which forms the foundation for many classical magnetic technologies.

Unlike FMs, AFs exhibit a more complex magnetic ordering characterized by the cancellation of magnetic moments at the macroscopic scale. This behavior can be effectively illustrated using a three-dimensional cubic unit cell, where the magnetic structure consists of two interpenetrating sublattices with moments aligned uniformly within each sublattice[†]. In a prototypical A-type AF, for instance, spins within individual atomic planes are ferromagnetically aligned,

[†]While the two-sublattice model captures basic antiferromagnetic behavior, many AFs exhibit more complex structures, some require multiple sublattices, and others feature non-collinear spin arrangements. A prominent example is the kagome AF, where spins form a 120° pattern across three sublattices.

while adjacent planes are coupled antiferromagnetically, resulting in alternating spin orientation along one crystallographic axis. This antiparallel configuration leads to complete cancellation of the overall magnetization, as illustrated in figure 1.1b. Although each sublattice carries a net magnetic moment, their vector sum is zero, resulting in no macroscopic magnetization and no stray magnetic fields. This absence of an external magnetic signature has historically made AFs significantly more challenging to detect and study using conventional magnetic measurement techniques. AFs, once seen as inactive materials, are now widely

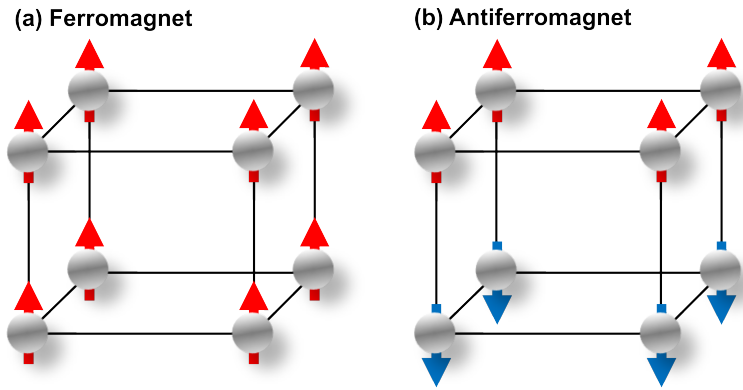


Figure 1.1: Schematic illustration of magnetic moment configurations within a cubic unit cell for (a) a ferromagnetic material, where all moments are aligned parallel, and (b) an A-type antiferromagnetic material, characterized by ferromagnetically aligned moments within each atomic plane and antiparallel alignment between adjacent planes. The red and blue arrows denote the two magnetic sublattices in the AF.

used in a passive role to pin FM layers via the exchange bias effect. When thin FM and AF films are coupled, the AF fixes the FM's orientation, remaining stable under moderate external fields^{9–11}. This passive use is key in magnetic sensors, such as hard disk drive (HDD)¹² read heads and magnetic random access memory (MRAM)^{13,14}, where one FM layer is pinned by an AF while the other responds to external magnetic or current induced effective fields.

Beyond this passive role, significant progress has recently been made in the active manipulation and detection of antiferromagnetic order. A central challenge is how to control and detect antiferromagnetic order. Achieving control with an external magnetic field generally demands fields of several tesla, and the lack of macroscopic magnetization makes antiferromagnetic order difficult to observe

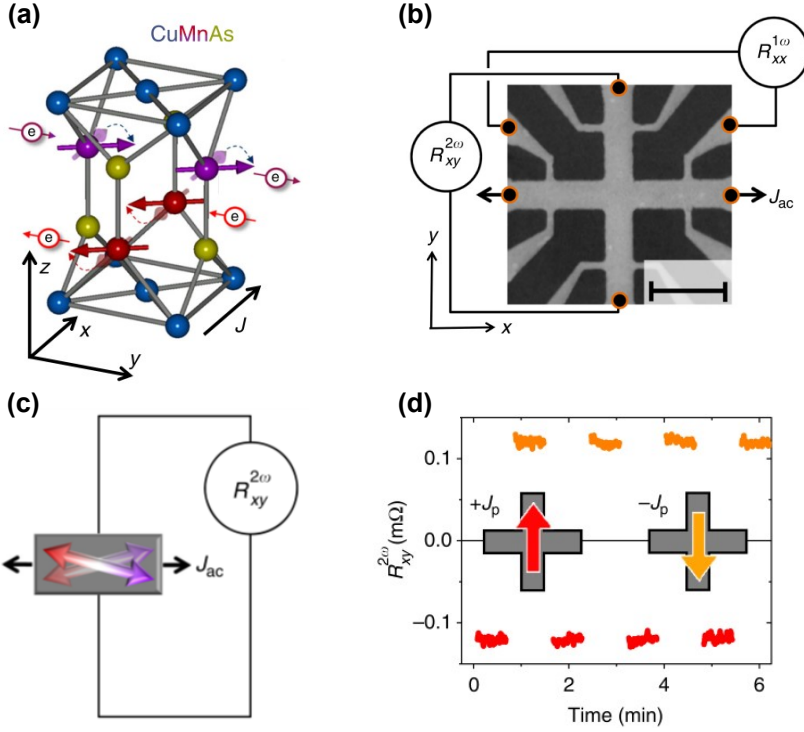


Figure 1.2: Electrical detection of 180° Néel vector reversal in CuMnAs. (a) A probing current (black arrow) generates staggered spin polarization (red/purple electrons) that deflects the antiferromagnetic moments (thick arrows on Mn sites). (b) Scanning electron micrograph of the cross-bar device with contacts for longitudinal ($R_{xx}^{1\omega}$) and transverse ($R_{xy}^{2\omega}$) resistance measurements; schematics illustrate the simultaneous detection setup. Scale bar: 10 μm . (c) Alternating probing current J_{ac} along x induces second-harmonic voltage along y , giving $R_{xy}^{2\omega}$; double arrows show the alternating deflection mechanism. (d) Writing pulses $J_p = 11$ mA ($j_p \sim 10^7$ A/cm 2) along $\pm y$ set the Néel vector along $\pm x$; $R_{xy}^{2\omega}$ is measured with $J_{ac} = 2$ mA ($j_{ac} \sim 10^6$ A/cm 2), pulse duration 20 ms. Reproduced from [15].

directly. The research group in which this doctoral research was conducted has played a foundational role in developing experimental and theoretical approaches for the electrical and optical manipulation and detection of antiferromagnetic order^{3,4,15–19}. Jörg Wunderlich, supervisor of the present thesis, was among the discoverers of the intrinsic spin Hall effect (SHE) – a phenomenon that underpins spin-orbit torque (SOT)-induced control of ferromagnetic and antiferromagnetic

order and forms a central mechanism explored in this thesis (more on this in [chapter II](#))²⁰. In materials with broken spatial inversion symmetry, relativistic spin-orbit coupling (SOC) gives rise to a current-induced staggered spin-orbit field^{16,20}, enabling efficient control of the antiferromagnetic order in AFs such as CuMnAs^{15,17,18,21} and Mn₂Au^{22–24}. In antiferromagnetic CuMnAs films, 180° reversals of the Néel vector can be controlled by switching the polarity of the writing current¹⁵. As shown in [figure 1.2](#), the two stable states with opposite Néel vector orientations in this collinear AF can be electrically distinguished via a second-order magnetoresistance effect enabled by SOTs. A key feature of antiferromagnetic memory devices is multilevel switching, enabling the integration of memory and logic within a single bit cell. This behavior, linked to multidomain reconfigurations, has been observed in metallic^{22,23}, semimetallic^{17,25}, and semiconducting²⁶ AF structures. While the intrinsic resonance frequencies of FMs are typically in the GHz range, AFs can exhibit tera hertz (THz)-range resonance due to strong inter-sublattice exchange, enabling much faster and more energy-efficient switching. Reversible switching via current-induced SOT has been demonstrated with pulse lengths reduced from seconds to picoseconds^{19,25}. While AFs can be actively controlled, they can also act as spin current sources to switch adjacent FMs, enabling field-free magnetization reversal in FMs²⁷.

Before electrical control of AFs, ultrafast optical studies had already demonstrated spin reorientation via both thermal and non-thermal excitations. Laser-induced heating and THz pulses altered anisotropy fields, enabling large-angle spin dynamics in ferrites^{28–30}. While most optical studies focus on insulating AFs, recent efforts are extending to semi-metallic and metallic materials used in electrical devices, bridging optical and electronic antiferromagnetic spintronics^{19,31}.

Recent studies extend beyond uniform antiferromagnetic order to explore domain walls (DWs) and complex spin textures like skyrmions. In crystalline AFs, DWs can be reversibly manipulated by an electrical current and visualized by X-ray magnetic linear dichroism¹⁸. Efficient DW and skyrmion motion has also been demonstrated in synthetic antiferromagnets (SAFs), making them promising platforms for racetrack memory technologies³².

Another promising direction lies at the intersection of spintronics and topology in AFs with complex non-collinear orders. Certain compensated non-collinear AFs with strong SOC exhibit large anomalous Hall effect (AHE)^{33,34} and anomalous Nernst effect (ANE)³⁵ despite having no net magnetization, due to the lack of

effective time-reversal symmetry and a finite Berry curvature. These non-collinear AFs can be manipulated electrically³⁶ and can also generate unusual transverse and longitudinal spin transport effects³⁷.

1.2 Synthetic antiferromagnets: a laboratory playground for real antiferromagnets

The rapid development of antiferromagnetic spintronics, driven by advances in electrical switching, ultrafast optical control, and topologically emergent transport phenomena, has highlighted both the promise and the practical challenges of working with crystalline AFs. Their inherently strong exchange fields, THz-scale dynamics, and lack of net magnetization endow them with remarkable functionalities but also render them difficult to probe, manipulate, and integrate into device architectures. In this context, SAFs have emerged as an invaluable model platform that captures the essential physics of antiferromagnetic order while remaining experimentally accessible^{38–43}. While SAFs differ from real AFs in microscopic exchange scales and dipolar contributions, they reproduce the key two-sublattice antiferromagnetic phenomena in an experimentally accessible regime.

SAFs are engineered multilayer structures in which two (or more) ferromagnetic layers are coupled antiferromagnetically across a metallic spacer[‡]. Although artificial, their compensated ground state, strongly suppressed stray fields, and exchange-enhanced collective modes parallel the behavior of real AFs. Yet, crucially, the inter-layer exchange coupling (IEC) in SAFs is orders of magnitude weaker than the intrinsic exchange in real AFs. This reduction in effective exchange field brings antiferromagnetic-like dynamics into the giga hertz (GHz) regime, making them compatible with conventional laboratory instrumentation and enabling systematic studies of phenomena that are otherwise hidden on ultrafast THz timescales.

Another key advantage of SAFs is tunability. SAFs allow precise control over exchange coupling, anisotropy, and magnetic configuration simply by adjusting layer thicknesses, material choice, or spacer composition. This tunability enables the stabilization of both in-plane and out-of-plane antiferromagnetic configura-

[‡]A thorough explanation of the SAFs can be found in [section 2.1.3](#)

tions, tailoring of domain behavior, and engineered access to distinct dynamical modes, including chiral excitations.

Owing to this combination of experimental accessibility, tunability, and representation of core antiferromagnetic physics, SAFs have become a foundational platform in modern spintronics research. They enable the exploration of DW dynamics, skyrmion motion, SOT-driven switching, and nonlinear magnetization dynamics, and they provide a versatile environment for investigating neuromorphic functionalities. These capabilities make SAFs not only a conceptual bridge between FMs and atomic AFs but also a practical testing ground for device-oriented antiferromagnetic spintronics.

For these reasons, this thesis adopts SAFs as its primary material platform, leveraging their controllability to study current-induced manipulation of antiferromagnetic order, investigate complex eigenmodes, and demonstrate dynamical phenomena such as chiral excitations and self-sustained oscillations. A comprehensive discussion of the IEC mechanism, magnetic configurations, and dynamical properties of SAFs is provided in the following [chapter II](#) and [chapter III](#).

1.3 Neuromorphic spintronics: opportunities and challenges

“In order to understand the brain, we have used the computer as a model for it. Perhaps it is time to reverse this reasoning. To understand where we should go with the computer, we should look to the brain for some clues”. This insight from Robert Noyce (founder of Intel Corporation) resonates more strongly today than ever before, a relevance underscored by the 2024 Nobel Prize in Physics awarded to John J. Hopfield and Geoffrey E. Hinton *“for foundational discoveries and inventions that enable machine learning with artificial neural networks”*⁴⁴. Artificial intelligence (AI) has become an integral part of modern life, from facial and pattern recognition to language translation, search engines, and the GPT-like large language models that now mediate vast amounts of human-computer interaction⁴⁵. While AI can already surpass human performance in several specialized tasks such as image or speech classification and strategic gaming, its practical expansion is hindered by a fundamental bottleneck: the limitations of the underlying comput-

ing hardware⁴⁶. Conventional processors rely on the von Neumann architecture, in which memory and logic units are physically separated. Although highly effective for precise numerical computation, this architecture is inherently inefficient for AI workloads, which must shuttle massive volumes of data back and forth between memory and processor^{47–49}. This “*von Neumann bottleneck*” results in high latency and substantial energy consumption. As shown in figure 1.3, recent analyses indicate that the global demand for computing power is now doubling approximately every two months, far outpacing the gains provided by Moore’s law. To sustain the growth of AI, new device concepts and alternative computing paradigms are urgently needed.

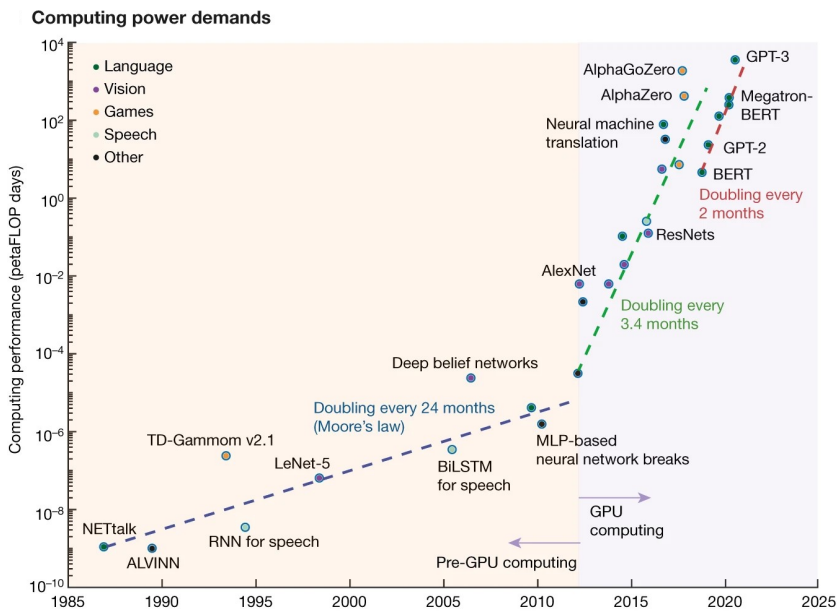


Figure 1.3: Growth of computing power demand over the past four decades, expressed in petaFLOPS days. Until 2012, the computational requirements of state-of-the-art applications doubled approximately every 24 months; in recent years, this doubling period has accelerated dramatically to roughly every 2 months. Reproduced from [47].

One promising direction is neuromorphic computing, which seeks to emulate the structure and working principles of the biological brain^{47–55}. Unlike modern processors, the brain performs computation using vast numbers of low-precision operations running in parallel, with memory stored locally in synaptic weights.

Neurons communicate using brief electrical spikes rather than continuous logic signals, enabling extraordinary computational efficiency. Emulating such behavior using conventional transistors, however, requires large and power-hungry circuitry, rendering it unsuitable for many embedded or battery-powered applications.

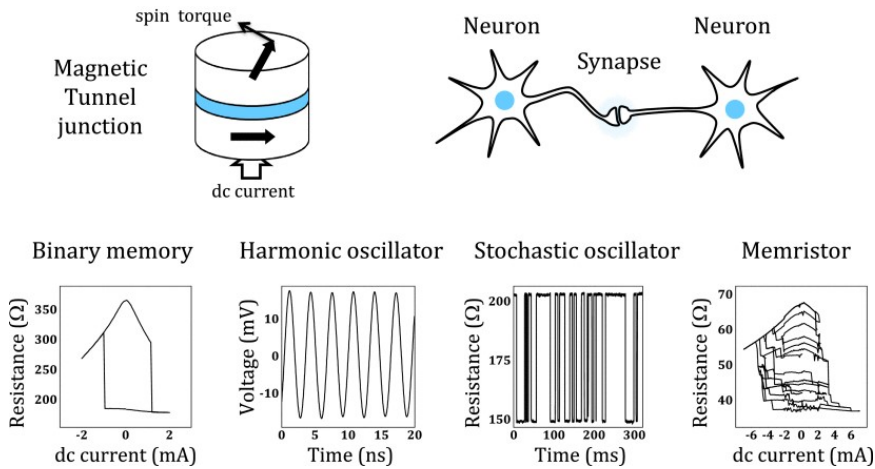


Figure 1.4: Principle and multifunctionality of a neuromorphic spintronic device. Top: A direct current (dc) injected through a magnetic tunnel junction, one representative implementation among several possible spintronic platforms, generates a spin torque acting on the magnetization of the active layer. The resulting magnetization dynamics modulate the device resistance, enabling the emulation of essential synaptic and neuronal functionalities. Bottom: By tailoring the device geometry and tuning external parameters such as applied magnetic field or drive current, a diverse set of dynamical responses can be achieved. Shown here are four representative modes: binary switching and memristive behavior, relevant for artificial synapses; and harmonic or stochastic oscillations, which replicate characteristic behaviors of biological neurons or neuronal assemblies. Reproduced from [56].

In this context, spintronic devices have emerged as compelling building blocks for brain-inspired hardware^{56–61}. These devices exploit the electron's spin degree of freedom to achieve functionalities, such as nonvolatility, analog tunability, plasticity, oscillatory dynamics, and intrinsic stochasticity, that naturally mirror key aspects of biological neurons and synapses⁶². Figure 1.4 illustrates the multifunctionality of a neuromorphic spintronic device, using a magnetic tunnel junction (MTJ) as a representative example platform. Importantly, these diverse

behaviors can often be realized using similar material stacks, enabling compact, multifunctional neuromorphic elements. Spintronic components can transmit information not only through charge currents, but also via spin currents, microwave signals, spin waves, DWs and topological magnetic textures such as skyrmions^{57,61}. This versatility opens pathways to neuromorphic architectures that are potentially far more energy-efficient, robust, and scalable than the transistor-based counterparts.

Analog, nonvolatile, and plastic resistors, commonly referred to as memristors, were first envisioned by Chua⁶³ and now represent a fundamental concept in neuromorphic computing. Most traditional memristive approaches for neuromorphic hardware, however, rely on ionic migration or phase transitions, which often suffer from limited reliability and endurance^{64–67}. Spintronic systems, by contrast, leverage magnetic phenomena that are inherently fast, durable, and compatible with high-density integration. Wang et al.⁶⁸ proposed a spintronic memristor based on the displacement of a magnetic DW within a spin-valve, where the device resistance varies depending on the DW position. Subsequently, Chanthbouala et al.⁶⁹ and Lequeux et al.⁷⁰ experimentally demonstrated memristive behavior in MTJs by exploiting controlled DW motion. Spintronic memristors based on topological magnetic textures, such as skyrmions, are also being actively explored for neuromorphic applications^{71,72}. However, a key limitation of these state-of-the-art ferromagnetic artificial synapses arises from long-range magnetic stray fields generated by the global magnetic configuration. These stray fields can perturb neighboring nanostructures, particularly during the writing process. Combined with thermal fluctuations and the probabilistic nature of magnetic moment dynamics, such interactions can lead to stochastic domain wall pinning and depinning, making it challenging to achieve reliable and reproducible writing of analog synaptic weights.

Spin transfer torques (STTs) and SOTs based nano-oscillators can be driven into spontaneous microwave oscillations by an injected direct current and possess several features highly appealing for neuromorphic computing^{73–84}. Their oscillation amplitudes exhibit memory due to finite magnetization relaxation, effectively emulating the leaky integration of biological neurons. Moreover, the frequency and amplitude of voltage oscillations are highly nonlinear functions of current or applied field, enabling the direct implementation of nonlinear activation functions. Coupling between spin-torque nano-oscillators can be achieved via direct

exchange interactions^{85–87}, magnetic fields^{88,89}, or oscillating electrical currents mediated by giant or tunneling magnetoresistance⁹⁰. Such interactions allow the formation of large arrays of coupled oscillators, emulating the synchronization of neurons and neural ensembles observed in the brain. For instance, Torrerón et al. demonstrated neuromorphic computation with a single spin-torque nano-oscillator by time-multiplexing to emulate a full network of 400 neurons⁹¹. Romera et al.⁹² used a short chain of four weakly coupled STT driven oscillators to perform vowel recognition, while Zahedinejad et al.^{93,94} recently reported two-dimensional synchronization in large arrays of SOT-based nano-oscillators, achieving robust operation of up to 64 strongly coupled devices at 10 GHz with a record-high quality factor.

Despite these advances, hitherto STT- and SOT-based oscillators and memristors rely on ferromagnetic materials, which remain sensitive to external magnetic fields that alter the energy landscape and device dynamics. Their operating speeds are limited to the GHz range, and their tunability, whether current-controlled frequency or chirality of oscillations, is constrained. Antiferromagnetic spintronics offers a promising solution^{3,4,61,95,96}: AF materials are inherently insensitive to even very large magnetic fields while providing ultrafast dynamics (in THz), complex and varied oscillation modes (more on this in [chapter III](#)), compactness, and other properties mentioned in [section 1.1](#) which are highly suitable for neuromorphic computing. Further leveraging antiferromagnetic spin textures, such as skyrmions, for neuromorphic spintronics can benefit from faster speed and their intrinsically suppressed skyrmion Hall effect (Magnus-force-driven transverse deflection) compared with ferromagnetic skyrmions. Moreover, the recently discovered class of altermagnetic materials, which exhibit antiferromagnetic ordering in real space while displaying ferromagnetic-like characteristics in electron momentum space, may provide new pathways to overcome these limitations⁹⁷.

1.4 Motivation and outline of the thesis and key experimental contributions

The rapid progress of antiferromagnetic spintronics and its growing relevance for neuromorphic technologies, as outlined in the previous sections, strongly motivates the need for experimentally accessible platforms that capture the es-

essential physics of AFs while enabling precise control and detection of their order. As has been detailed in [section 1.2](#), SAFs provide such a model system. In this thesis, SAFs are therefore used as a versatile laboratory platform to study current induced SOT-driven manipulation, dynamics, and oscillations of the Néel order, phenomena that are central to both antiferromagnetic spintronics and its emerging neuromorphic applications.

To establish the necessary conceptual foundation, the thesis first introduces (in [chapter II](#) and [chapter III](#)) the static and dynamical aspects of antiferromagnetic spintronics. [Chapter II](#) surveys the microscopic interactions governing antiferromagnetic order, the structure of antiferromagnetic domains and DWs, and the SOT mechanisms, rooted in the inverse spin galvanic effect (iSGE) and the SHE, that enable electrical control of the Néel order. It further outlines the magnetotransport and thermoelectric responses that provide electrical access to the Néel vector. [Chapter III](#) turns to magnetization dynamics, beginning with the Landau-Lifshitz-Gilbert (LLG) equation and ferromagnetic resonance (FMR), and then developing a coupled-sublattice description of antiferromagnetic eigenmodes and spin waves (SWs). The chapter concludes with the discussion on SOT-driven self-oscillations, a key phenomenon explored in this thesis.

[Chapter IV](#), [chapter V](#) and [chapter VI](#) present the three main experimental results. Each addresses a distinct aspect of SAF-based antiferromagnetic functionality, ranging from SOT-driven nonvolatile switching of Néel order to SOT enabled time-resolved magneto-optical Kerr effect (tr-MOKE) imaging of antiferromagnetic domains and their dynamics and the realization of a nanoscale SOT-driven self-oscillator. A brief overview is given below.

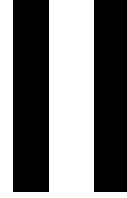
[Chapter IV](#) demonstrates electrical writing and readout of the Néel order in a parity-time ($\mathcal{PT}$)-symmetric perpendicular-anisotropy SAF. SOT drives controlled DW motion, enabling stable and reconfigurable nonvolatile states, while nonlinear magnetotransport provides a sensitive electrical probe of the underlying antiferromagnetic texture. The absence of long-range stray fields allows reproducible access to intermediate domain configurations, establishing the SAF as a compact, robust spintronic memristor capable of encoding analog synaptic weights.

[Chapter V](#) investigates current-induced antiferromagnetic dynamics in an in-plane easy-plane SAF using stroboscopic time-resolved magneto-optical Kerr

effect (tr-MOKE) microscopy. A polar magneto-optical Kerr effect (p-MOKE) geometry, normally insensitive to in-plane SAF antiferromagnetic order, is rendered sensitive by a staggered damping-like spin-orbit field that acts oppositely on the two magnetic layers, generating a dynamic out-of-plane magnetic moment and enabling real-space imaging of antiferromagnetic domains. Phase-resolved measurements show that the Néel order in adjacent domains oscillates 180° out of phase, providing direct evidence of antiferromagnetic domain dynamics. In addition, DC currents drive reversible DW motion, which is observed in these tr-MOKE measurements. Finally, the chapter demonstrates a strongly localized antiferromagnetic excitation confined to the constriction, highlighting the role of device geometry in shaping antiferromagnetic dynamics.

The final experimental chapter ([chapter VI](#)) reports the realization of a SOT-driven nano-oscillator in an easy-plane SAF nanoconstriction. Spin rectification measurements identify current-excited optical and acoustic eigen modes, as well as additional resonances emerging above a critical DC current near the spin-flop transition. These resonances correspond to self-sustained chiral oscillations whose handedness is determined by current polarity and which can be injection-locked to an external microwave drive. Complementary macrospin and micromagnetic simulations confirm the existence of these oscillations and predict the onset of chaotic dynamics in the vicinity of the spin-flop transition, consistent with the experimentally observed multi-mode behavior.

Finally, [chapter VII](#) synthesizes the main findings of the thesis, situates them within the broader landscape of antiferromagnetic spintronics, and outlines opportunities for future research, including pathways toward implementing SAF-based neuromorphic elements and extending the demonstrated concepts to real AFs and altermagnets.



Fundamentals of antiferromagnetic spintronics

Following the brief introduction of antiferromagnetic spintronics in [chapter 1](#), this chapter develops the theoretical and phenomenological foundations relevant to the static properties of AFs. In [section 2.1](#), we outline the fundamental interactions relevant to the antiferromagnetic order with an emphasis on SAF that constitutes the primary material system studied in this thesis. [Section 2.2](#) addresses the formation and stability of antiferromagnetic domains and DWs. [Section 2.3](#) introduces the essential spintronic concepts required for the manipulation of antiferromagnetic order, including spin currents, the SHE, and SOTs. Particular attention is given to SOTs, as it represents the sole mechanism employed in this thesis for electrical control and excitation of antiferromagnetic order and its dynamics. Finally, [section 2.4](#) and [section 2.5](#) review the magnetotransport and thermoelectric effects used to detect the antiferromagnetic state. Together, these sections establish the theoretical framework underpinning the experimental and dynamical investigations presented in subsequent chapters.

Keywords: antiferromagnetic spintronics, synthetic antiferromagnets (SAFs), spin Hall effect (SHE), spin-orbit torques (SOTs), antiferromagnetic domains, magnetoresistances.

2.1 Magnetic interactions in antiferromagnets

Magnetic ordering originates from quantum mechanical interactions. While a rigorous understanding of magnetism at the microscopic level requires a quantum mechanical treatment on the atomic level, such a detailed description becomes unnecessary when considering macroscopic properties that emerge on length scales much larger than interatomic distances. At low energies, the behavior of magnetic materials is dominated by collective interactions between the individual atomic spins. In this context, the following will focus on a phenomenological description of the key interactions governing antiferromagnetic systems.

2.1.1 Antiferromagnetic ordering

The equilibrium properties of FMs are explained through Pierre Weiss's concept of a global molecular field¹. Building on the Curie law for paramagnets (PMs), where the susceptibility satisfies $\chi^{-1} \approx T$, Weiss proposed that FMs experience an additional internal field $\lambda \mathbf{M}$, proportional to the magnetization $\mathbf{M}$. This leads to the Curie-Weiss law at high temperatures: $\chi^{-1} \approx T - \theta$, where $\theta > 0$ is the Curie-Weiss temperature and is proportional to the Weiss constant λ ^{7,8}.

In the early 1930s, Néel addressed the puzzling behavior of certain magnetic materials that exhibited zero remanence at all temperatures yet deviated from the paramagnetic Curie law¹. These materials followed the Curie-Weiss law at high temperatures with a *negative* Curie-Weiss temperature θ , and displayed nearly constant susceptibility at low temperatures. Using a model of two interlaced cubic sublattices, Néel proposed a new type of magnetic order, *antiferromagnetism*, in which the local molecular fields act in opposite directions on each sublattice, resulting in equal and opposite spontaneous magnetizations that cancel out (figure 1.1b).

In magnetically isotropic systems (i.e., neglecting spin-lattice coupling), a weak external magnetic field favors a configuration where the Néel vector lies perpendicular to the field, with sublattice moments canting to produce a small net magnetization. As the field strength increases, the sublattice moments cant toward the field direction, producing a net magnetization that is proportional to the applied field, inversely proportional to the local molecular field constant,

and independent of temperature. This mechanism accounts for the constant low-temperature susceptibility observed in many AFs.

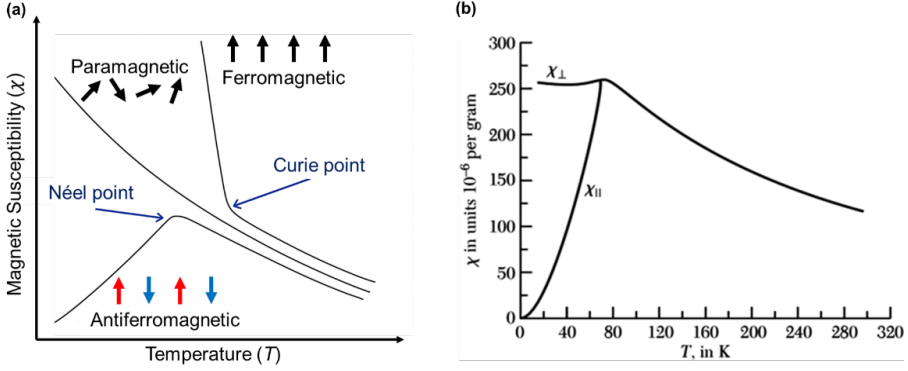


Figure 2.1: Representative magnetic susceptibility versus temperature for paramagnetic, ferromagnetic, and antiferromagnetic materials (a), and experimental magnetic susceptibility data for MnF_2 (b) (reproduced from [98]).

Figure 2.1a shows the temperature dependence of susceptibility for PMs, FMs, and AFs, while figure 2.1b presents the measured parallel and perpendicular susceptibility components of antiferromagnetic MnF_2 . For FMs, the natural order parameter is the net magnetization $\mathbf{M}$, defined as the sum of all magnetic moments in the material. In AFs, however, where moments prefer antiparallel alignment, the spontaneous net magnetization $\mathbf{M}$ vanishes at zero magnetic field^{*}, necessitating a more appropriate order parameter. A simple model divides the AF structure into two interpenetrating sublattices with opposing magnetizations $\mathbf{M}_1 = -\mathbf{M}_2$.

Rather than treating the sublattices separately, it is convenient to define the *total magnetization* $\mathbf{M} = (\mathbf{M}_1 + \mathbf{M}_2)/2M_s$, and the *staggered magnetization* or *Néel vector* $\mathbf{L} = (\mathbf{M}_1 - \mathbf{M}_2)/2M_s$, where M_s is the saturation magnetization of each sublattice. The vector $\mathbf{L}$ characterizes the antiferromagnetic spin alignment and serves as the order parameter for collinear AF order. Perfect antiferromagnetic order corresponds to $|\mathbf{L}| = 1$, $|\mathbf{M}| = 0$, while perfect ferromagnetic order

^{*}The net magnetization $\mathbf{M}$ vanishes only in the absence of an external magnetic field. However, in certain materials known as *weak ferromagnets*, intrinsic interactions such as the Dzyaloshinskii-Moriya interaction can give rise to a small but finite net magnetization. Examples include $\alpha\text{-Fe}_2\text{O}_3$, MnCO_3 , and CoCO_3 .

corresponds to $|\mathbf{L}| = 0$, $|\mathbf{M}| = 1$. Although neither $\mathbf{L}$ nor $\mathbf{M}$ are unit vectors, they satisfy the constraints $\mathbf{L}^2 + \mathbf{M}^2 = 1$ and $\mathbf{L} \cdot \mathbf{M} = 0$.

2.1.2 Exchange interaction

The interaction responsible for magnetic ordering in solids is the *exchange interaction*, which is a purely quantum mechanical effect and originates from the combination of Coulomb repulsion between nearby electrons and the Pauli exclusion principle, that prevents two electrons from occupying the same quantum state^{7,8}. Because the total electronic wavefunction must be antisymmetric under particle exchange, the spatial symmetry of a two-electron state is directly linked to its spin configuration. This symmetry constraint leads to an energy difference between spin singlet ($S = 0$) and triplet ($S = 1$) states, even in the absence of any explicit spin-dependent term in the Coulomb interaction⁸.

For two interacting spins $\mathbf{S}_1$ and $\mathbf{S}_2$, this energy difference can be expressed in terms of the scalar product $\mathbf{S}_1 \cdot \mathbf{S}_2$, leading to the effective spin dependent Hamiltonian

$$\mathcal{H}^{\text{spin}} = -2J \mathbf{S}_1 \cdot \mathbf{S}_2, \quad (2.1)$$

where J is the exchange constant (exchange integral). The sign of J determines the preferred spin alignment: $J > 0$ favors parallel alignment (ferromagnetic coupling), while $J < 0$ favors antiparallel alignment (antiferromagnetic coupling). This simplified picture captures the essence of the exchange interaction at the microscopic level.

In magnetic materials, similar exchange couplings occur between localized spins on neighboring atoms, described by the Heisenberg Hamiltonian

$$\mathcal{H}_E = -2 \sum_{i>j} J_{ij} \mathbf{S}_i \cdot \mathbf{S}_j, \quad (2.2)$$

where J_{ij} is an isotropic exchange parameter that quantifies the coupling between classical spins $\mathbf{S}_i$ and $\mathbf{S}_j$ located at lattice sites i and j . Since J_{ij} depends on the overlap of atomic orbitals, the exchange interaction is typically short-ranged and often restricted to nearest neighbors, allowing the approximation $J_{ij} \rightarrow J$. A negative J indicates antiferromagnetic coupling whereas a positive J corresponds

to ferromagnetic interactions. In both FMs and AFs, the exchange interaction commonly dominates spin behavior.

2.1.3 RKKY interaction and synthetic antiferromagnets

In addition to the direct exchange interaction discussed in [section 2.1.2](#), metallic systems also exhibit an *indirect exchange interaction* mediated by conduction electrons. In this mechanism, a localized magnetic moment induces spin polarization in the surrounding conduction electrons, which in turn mediates an effective coupling to another magnetic moment at a distance r . This *itinerant exchange*, known as the *Ruderman-Kittel-Kasuya-Yosida (RKKY) interaction*, originates from the oscillatory spin response of the conduction electrons^{7,8,38–41,99,100}. The interaction strength is given by:

$$J_{\text{RKKY}}(r) \propto \frac{\cos(2k_F r)}{r^3}, \quad (2.3)$$

where k_F denotes the radius of the spherical Fermi surface. The RKKY interaction is *long-range* and *oscillatory*, and depending on the separation r , it can favor either ferromagnetic or antiferromagnetic alignment of the magnetic moments[†].

Recent advances in nanofabrication have made it possible to realize *nanometer-thin multilayer structures* composed of alternating magnetic and non-magnetic (NM) metallic layers. In such systems, interlayer RKKY coupling arises from spin dependent Friedel-like oscillations in the spin density of the NM spacer, induced by the adjacent FMs^{99,101–107}. By tuning the spacer thickness, the coupling can be controlled, favoring *parallel (ferromagnetic)* alignment at certain thicknesses and *antiparallel (antiferromagnetic)* alignment at others ([figure 2.2](#)). For sufficiently thick spacers, the interaction is suppressed. Multilayer structures in which two ferromagnetic layers are antiferromagnetically coupled via this mechanism, separated by a NM spacer, are known as *synthetic antiferromagnets (SAFs)* (see [figure 2.3](#)). More complex structures such as SAF-based multilayers or superlattices

[†]The form of IEC considered in this thesis is the *bilinear exchange*³⁹, expressed as $-J \mathbf{M}_1 \cdot \mathbf{M}_2$. In addition to this, a higher-order term known as *biquadratic exchange*³⁹ can also arise, described by $-J_2 (\mathbf{M}_1 \cdot \mathbf{M}_2)^2$. Experimental studies have shown that the biquadratic coupling constant J_2 is typically negative, which energetically favors a perpendicular (90°) orientation between the two magnetizations. However, in this thesis, we focus exclusively on bilinear exchange coupling that favors antiparallel alignment between the magnetizations of the ferromagnetic layers.

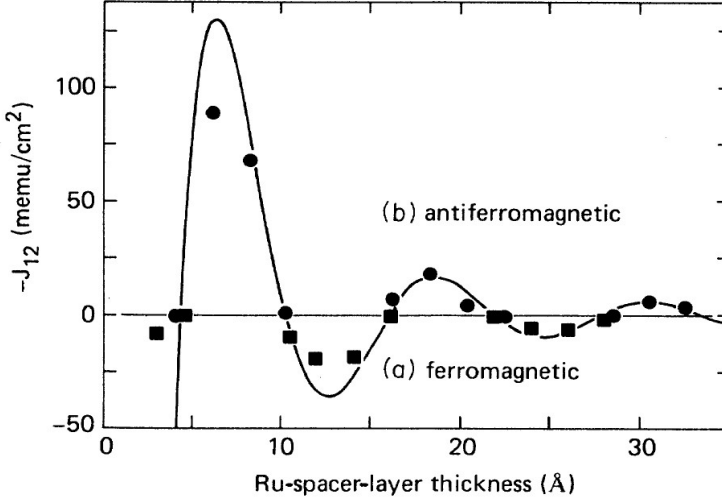


Figure 2.2: Interlayer RKKY exchange coupling strength J_{12} as a function of Ru spacer layer thickness for coupling between $\text{Ni}_{80}\text{Fe}_{20}$ layers. The oscillatory nature of the coupling demonstrates alternating ferromagnetic and antiferromagnetic alignment as the spacer thickness varies (reproduced from [106]).

can also be realized, in which each alternate ferromagnetic layer is antiferromagnetically coupled to its adjacent ferromagnetic layers through RKKY-mediated interlayer exchange⁴¹.

This antiferromagnetic coupling between magnetic layers in SAFs was instrumental in uncovering the dependence of electrical resistance in metallic multilayers on the relative alignment of adjacent magnetizations^{108,109}. This effect, termed giant magnetoresistance (GMR), or tunnel magnetoresistance (TMR) when involving an insulating barrier, led the Nobel prize of 2007 and the foundation for the development of spintronics as a distinct research field¹¹⁰.

In comparison to the natural AFs with strong exchange interaction coupling strength (typically several 10-100 teslas), SAFs have much smaller IEC, thus enabling easier control of their antiferromagnetic order. Additionally, the magnetic state in SAFs can be readily probed through conventional techniques such as the magneto-optical Kerr effect (MOKE) or AHE (especially when the ferromagnetic layers differ slightly in their properties)⁴¹. Consequently, while their magnetiza-

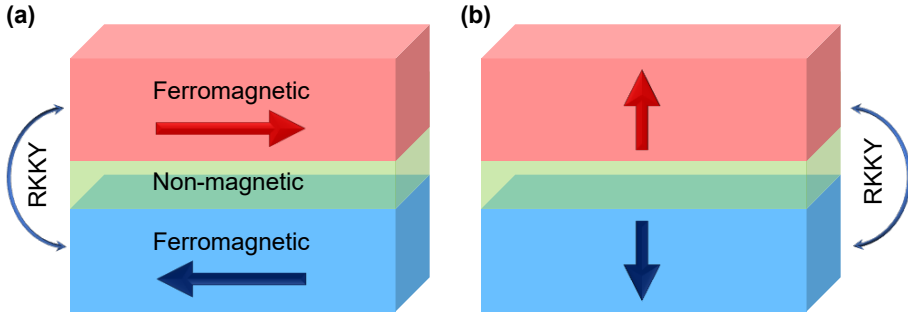


Figure 2.3: Schematic of synthetic antiferromagnets. (a) Bilayers with in-plane magnetization. (b) Bilayers with out-of-plane magnetization. The arrows within each ferromagnetic layer indicate the direction of magnetization (adapted from [41]).

tion dynamics exhibit similarly complex behavior to natural AFs, SAFs operate in the GHz frequency range, making them far more accessible than natural AFs whose dynamics typically occur at THz frequencies.

SAFs offer exceptional tunability through anisotropy engineering, enabling the realization of both in-plane and out-of-plane magnetized configurations (see [figure 2.3](#)). The strength of the IEC can also be precisely controlled by adjusting the composition and thickness of individual ferromagnetic layers as well as the NM layer. Moreover, the magnetic order in SAFs can be efficiently manipulated via SOT mechanisms that exploit the SHE, achieved by placing heavy metal (HM) layers adjacent to the SAF. Despite being artificial structures, SAFs serve as excellent model systems for natural AFs, retaining the essential physics necessary to investigate antiferromagnetic behavior. Leveraging these advantages, this thesis focuses on current-induced manipulation of antiferromagnetic ordering in SAFs and their electrical detection, aiming to realize a stable, nonvolatile analog memristor. Furthermore, we probe the complex eigenmodes of SAFs, including chiral modes, and, to our knowledge for the first time, demonstrate DC current-induced self-oscillations, thereby realizing an antiferromagnetic oscillator. The magnetization dynamics are explored using both electrical and optical measurement techniques as shall be discussed in the upcoming chapters.

2.1.4 Demagnetizing field and shape anisotropy

In FMs, the *demagnetizing field* $\mathbf{H}_d$ (*stray field* outside the magnet) arises as a consequence of Maxwell's equations applied to a magnetized body, enforcing continuity of the magnetic induction^{7,8};

$$\mathbf{B} = \mu_0 (\mathbf{H} + \mathbf{M}), \quad (2.4)$$

where $\mathbf{B}$ is the total magnetic induction and $\mathbf{H}$ is the magnetic field. The distribution of $\mathbf{H}$, $\mathbf{M}$ and $\mathbf{B}$ is shown in [figure 2.4](#). Whenever a finite magnetization $\mathbf{M}$ exists, magnetic “charges” appear at surfaces (from $\nabla \cdot \mathbf{M}$) and at interfaces, producing an internal magnetic field that opposes the magnetization direction^{7,8}. This field is strongly dependent on the geometry of the sample and the boundary conditions at its surfaces. For ellipsoids, the demagnetizing field can be expressed analytically as

$$\mathbf{H}_d = -\hat{N} \cdot \mathbf{M}, \quad (2.5)$$

where $\hat{N}$ is the demagnetization tensor, generally represented by a symmetric 3×3 matrix determined only by the geometry^{7,8}. As $\mathbf{H}_d$ and $\mathbf{M}$ are collinear along the principal axes of the ellipsoid, the principal components of $\hat{N}$ in diagonal form (N_x, N_y, N_z) are known as demagnetizing factors. Demagnetizing factors are commonly used to obtain approximate internal fields, even for non-ellipsoidal shapes or numerical micromagnetic solutions are typically required⁷. The associated demagnetizing energy is

$$E_d = -\frac{1}{2}\mu_0 \int \mathbf{M} \cdot \mathbf{H}_d dV, \quad (2.6)$$

which is minimized when the magnetization lies along directions with the smallest demagnetizing factor. This produces the so-called *shape anisotropy*. For example, in thin films the out-of-plane demagnetizing factor approaches unity, while the in-plane components are nearly zero. As a result, thin ferromagnetic films exhibit a strong easy-plane anisotropy, favoring magnetization within the film plane unless overcome by perpendicular magnetic anisotropy (PMA)^{III}. The demagnetizing energy associated with the magnetic charges can be very large and thus unwanted which is referred to as the *charge avoidance principle*^{7,8,III}. This leads to a breakup of the uniform magnetization to magnetic domains which will be discussed later.

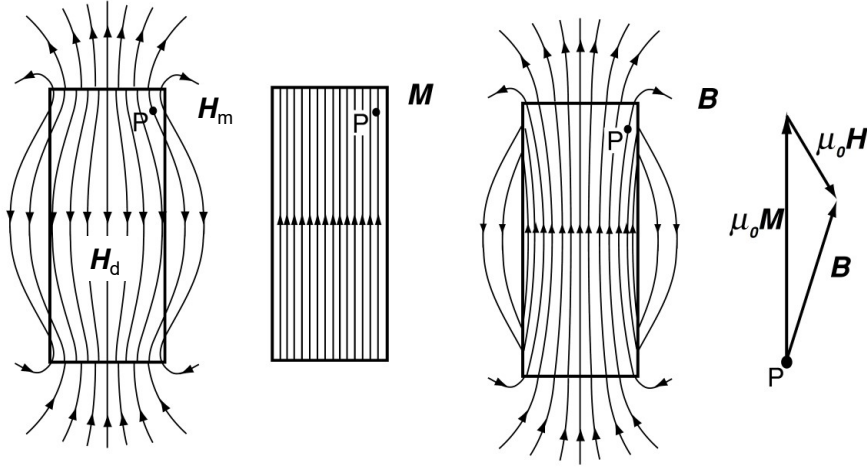


Figure 2.4: Distribution of magnetic fields generated by a uniformly magnetized body. Field lines of $\mathbf{B}$ are closed loops whereas the stray field ($\mathbf{H}_m$) emanates from the fictitious charges associated with the discontinuous magnetization at the top and bottom surface. The relation between the three vectors is shown in point P inside the magnet where $\mathbf{H}_m$ pointing against $\mathbf{M}$ and $\mathbf{B}$ is called demagnetizing field $\mathbf{H}_d$ (reproduced from [7]).

In an ideal, perfectly compensated collinear AF, the magnetic sublattices carry equal but opposite moments, such that the macroscopic magnetization vanishes. Consequently, the demagnetizing field ($\mathbf{H}_d$) and hence the conventional shape anisotropy is absent in perfectly compensated AFs⁵. However, antiferromagnetic domains are ubiquitous and their origin is not magnetostatic, but is governed by other long-range interactions which will be discussed in later section regarding antiferromagnetic domains.

In SAFs, while the ground state corresponds to antiparallel alignment of the layer magnetizations, perfect compensation is rarely achieved in practice due to unequal layer thicknesses, saturation magnetizations, or local nonuniformities. Each ferromagnetic layer continues to generate its own demagnetizing field and shape anisotropy, which owing to the antiparallel alignment results in a reduction of the stray fields. In addition, dipolar (magnetostatic) coupling between the layers persists: stray fields at the edges and surfaces of each layer interact with the other, modifying the total anisotropy landscape and the stability of domains⁴¹. In thin-film SAFs, the intra-layer shape anisotropy enforces an easy-plane anisotropy, such that both layer magnetizations (and hence the Néel vector) lie in-plane

unless a strong PMA is introduced. This has direct consequences for DW type (typically Néel-like in in-plane SAFs) and their dynamics under SOTs. In practical devices, geometry plays a central role. Nanoconstrictions (one of the studied device types in this thesis) concentrate magnetostatic charges at the constricted edges. In SAF nanoconstriction devices, this produces locally enhanced stray fields, altering the nucleation and motion of DWs as well as localization of spin wave (magnon) modes.

2.1.5 Magnetic anisotropy

While the exchange interaction enforces antiparallel alignment between neighboring spins in an AF, it is inherently isotropic, and it imposes no preference for an absolute spatial direction, only that the spins are oppositely aligned relative to each other. Magnetic anisotropy breaks this degeneracy and refers to the directional dependence of a material's magnetic properties, where magnetic moments preferentially align along specific crystallographic directions, known as easy axes^{7,8}. This behavior originates from the symmetry and atomic arrangement of the crystal lattice and is present in both FMs and AFs. Depending on the symmetry of the anisotropy, there can be one easy axis (uniaxial anisotropy), or a combination of multiple ones (e.g. two in case of a biaxial anisotropy)⁷.

We distinguish between *intrinsic magnetocrystalline anisotropy*, stemming from the ideal crystal symmetry, and *extrinsic anisotropies*, which arise from strain, surfaces, or interfaces. Magnetic anisotropy is temperature-dependent and vanishes at the magnetic ordering temperature (Curie temperature T_C or Néel temperature T_N)⁷. Different anisotropy contributions typically exhibit distinct temperature dependencies, enabling control over easy axis (EA) orientation and magnetization dynamics. Magnetocrystalline anisotropy arises from two main contributions:

- **Single-ion anisotropy**, driven by the interaction between the magnetic ion's electron cloud and the surrounding crystal field. SOC aligns the magnetic moment along preferred crystallographic directions, determined by the orbital symmetry. This interaction can be described using a multipole expansion involving spherical harmonics to capture the angular dependence of the anisotropy energy⁷.

- **Two-ion anisotropy**, originating from anisotropic dipole-dipole interactions between magnetic moments. The dipolar energy depends on the relative orientation of magnetic moments and the crystal structure, and may be significant in non-cubic lattices where the dipolar tensor does not vanish⁷.

The total anisotropy energy is often described by leading terms for cubic (E_c) and uniaxial (E_u) systems:

$$E_c = \frac{K_c}{4} (\sin^2 2\phi \sin^4 \theta + \sin^2 2\theta), \quad (2.7)$$

$$E_u = K_u \sin^2 \vartheta, \quad (2.8)$$

where θ and ϕ are the polar and azimuthal angles of the magnetization relative to the crystal axes, and ϑ is the angle between the magnetization and the easy axis in a uniaxial system. The anisotropy constants K_c and K_u determine the preferred directions. The corresponding anisotropy fields are defined as:

$$H_a = \frac{2K_a}{\mu_0 M_s}, \quad a = c, u, \quad (2.9)$$

where M_s is the saturation magnetization and μ_0 is the vacuum permeability.

In thin films, **induced anisotropies** become significant due to lattice mismatch with the substrate, surface symmetry breaking, or strain. These effects can be engineered, for example, via piezoelectric actuators or nanopatterning. Strain modifies both single-ion and two-ion contributions by distorting the crystal field and altering interatomic distances. Surface effects are especially important in ultrathin layers, where reduced atomic coordination and symmetry strongly influence magnetic behavior.

2.1.6 Dzyaloshinskii-Moriya interaction

The Dzyaloshinskii-Moriya interaction (DMI) is an antisymmetric exchange coupling that favors noncollinear spin arrangements with a fixed chirality. First proposed by Dzyaloshinskii in 1958 to explain weak ferromagnetism in certain AFs¹¹³, its microscopic origin was later attributed by Moriya (1960) to SOC combined

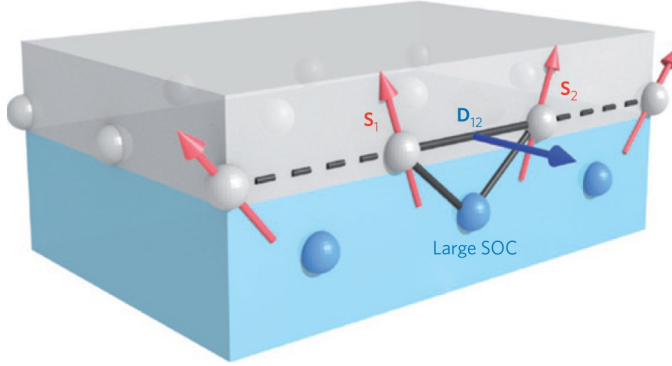


Figure 2.5: Schematic illustration of interfacial Dzyaloshinskii-Moriya interaction (DMI) at the interface between a ferromagnetic metal (grey) and a heavy metal layer with strong spin-orbit coupling (blue). The DMI vector $\mathbf{D}_{12}$ is associated with the triangle formed by two neighboring magnetic atoms and a heavy atom at the interface. It points perpendicular to the plane of the triangle, consistent with Moriya's symmetry rules. Reproduced from [112].

with broken inversion symmetry¹¹⁴. The DMI is described by the Hamiltonian:

$$\mathcal{H}_{\text{DMI}} = \sum_{\langle ij \rangle} \mathbf{D}_{ij} \cdot (\mathbf{S}_i \times \mathbf{S}_j), \quad (2.10)$$

where $\mathbf{D}_{ij}$ is the DMI vector, and $\langle ij \rangle$ denotes summation over nearest-neighbor spin pairs. The antisymmetric nature of the interaction favors a specific rotational sense between neighboring spins and competes with the symmetric Heisenberg exchange, although it is typically weaker in magnitude.

Despite its relative weakness, DMI plays a critical role in stabilizing a range of chiral magnetic textures, such as skyrmions and helical spin structures, which are widely observed in FMs and have been theoretically predicted in antiferromagnets^{96,112}. It also contributes to the stabilization of fixed-chirality ground states in certain noncollinear antiferromagnets^{33,34}.

Two types of DMI are generally distinguished: *bulk DMI*, which arises in crystals lacking inversion symmetry, and *interfacial DMI*, which occurs at interfaces where strong SOC and broken inversion symmetry are present. In this thesis, the studied SAFs are sandwiched between two HM layers, resulting in a net interfacial DMI due to the asymmetric interfaces. This interaction enforces specific DW chir-

alities and leads to the formation of *canted DWs* carrying a net magnetic charge, significantly affecting DW dynamics and stability (see [chapter IV](#)). [Figure 2.5](#) shows a schematic representation of the mechanism of the interfacial DMI in a FM/HM heterostructure.

2.1.7 Zeeman interaction and spin-flop/spin-flip transition

Apart from intrinsic interactions that determine the orientation of magnetic moments, such as exchange coupling and magnetic anisotropy, AF systems are also influenced by external magnetic fields. The interaction between the magnetic field $\mathbf{H}$ and magnetic moments is described by the Zeeman effect^{7,8,115}. On the atomic scale, this is captured by the Zeeman Hamiltonian,

$$\mathcal{H}_Z = - \sum_i g \mu_B \mathbf{S}_i \cdot \mathbf{H}, \quad (2.11)$$

where g is the Landé g -factor, μ_B is the Bohr magneton, and $\mathbf{S}_i$ is the spin operator for the i -th atom. At the macroscopic level, the Zeeman energy density for a magnetized body is given by

$$E_Z = -\mu_0 \mathbf{M} \cdot \mathbf{H}, \quad (2.12)$$

where $\mathbf{M}$ is the magnetization and μ_0 is the vacuum permeability. In a collinear AF, the two magnetic sublattices $\mathbf{M}_1$ and $\mathbf{M}_2$ are aligned antiparallel, resulting in a net magnetization $\mathbf{M} = \mathbf{M}_1 + \mathbf{M}_2 = 0$ in zero applied magnetic field. At first glance, this suggests that AFs are unaffected by external magnetic fields. Indeed, for moderate fields, the net moment remains essentially zero. However, a more accurate picture emerges by considering the effect of $\mathbf{H}$ on each sublattice separately.

Because the Zeeman term acts on each sublattice independently, the field tries to cant both $\mathbf{M}_1$ and $\mathbf{M}_2$ toward its direction. Due to their antiparallel coupling via the exchange interaction, this canting is constrained: when $\mathbf{M}_1$ cants slightly toward the field, $\mathbf{M}_2$ cants slightly away, maintaining a nearly antiparallel configuration and leaving the Néel vector unperturbed. This constraint can be overcome in two distinct ways in the case of an easy-axis AF:

When the external field is applied *perpendicular to the easy axis* (i.e., perpendicular to the Néel vector), energy can be reduced by canting the sublattices. This results in a small net magnetization along $\mathbf{H}$, while the Néel vector $\mathbf{L}$ shrinks in magnitude but retains its orientation. The exchange energy increases, but the Zeeman energy decreases more, favoring this configuration. This remains valid up to a critical field, beyond which the system undergoes a *spin-flip* transition where the sublattices align in the same direction and along the external field $\mathbf{H}$, resulting in a ferromagnetic-like configuration. The spin-flip field which is strong enough to fully overcome the anisotropy and exchange coupling, is then defined as¹¹⁵

$$\mu_0 H_{\text{spin-flip}} = 2H_E, \quad (2.13)$$

where we have assumed that the strength of exchange field $H_E \gg H_A$. H_A is the strength of the anisotropy field.

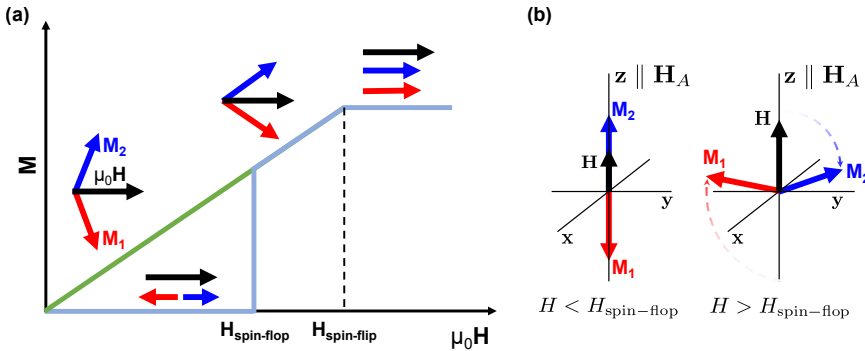


Figure 2.6: (a) Net magnetization and response of an antiferromagnet as a function of external magnetic field strength. The black arrows indicate the direction of the applied magnetic field $\mathbf{H}$, while the red and blue arrows represent the orientations of the sublattice magnetizations. Two scenarios are illustrated: when the initial Néel vector is perpendicular to the field (green curve), the sublattice moments gradually cant toward the field direction, resulting in a smooth increase of net magnetization; when the Néel vector is initially aligned with the field (cyan curve), the sublattice moments remain rigid until the spin-flop field $H_{\text{spin-flop}}$ is reached. At this point, a spin-flop transition occurs, reorienting the moments perpendicular to the field with canting that generates a net magnetization. Upon further increase of the field, the spin-flip transition occurs, leading to both sublattice magnetizations aligning parallel to the external field. This final state is reached for both initial configurations. (b) Illustration of the spin-flop transition.

When $\mathbf{H}$ is applied *along the easy axis* (i.e., parallel to the Néel vector), the sys-

tem faces a competition between the Zeeman energy and the magnetic anisotropy. Initially, the sublattice moments remain antiparallel along the easy axis. As the field increases, canting becomes energetically favorable, but the anisotropy acts as an energy barrier preventing immediate rotation of the Néel vector. At a critical field strength, called the *spin-flop field* ($H_{\text{spin-flop}}$), the Néel vector undergoes a sudden reorientation to lie perpendicular to $\mathbf{H}$, accompanied by symmetric canting of the sublattices. The spin-flop field can be estimated as¹¹⁵

$$\mu_0 H_{\text{spin-flop}} = 2\sqrt{H_A H_E}. \quad (2.14)$$

Figure 2.6 shows a schematic representation of the spin-flop and spin-flip phenomena in an easy-axis AF, illustrating the response of the sublattice magnetizations to an external magnetic field applied both along or perpendicular to the easy axis.

In natural AFs, the exchange field H_E can be extremely large (typically ~ 100 – 1000 T), which leads to very high spin-flop fields ($\mu_0 H_{\text{spin-flop}} \sim 10$ – 50 T)^{7,115}. These values are often impractical for applications and difficult to achieve experimentally. In contrast, SAFs exhibit significantly lower interlayer exchange (in the range of few mT to a few hundreds of mT). As a result, spin-flop and spin-flip transitions can occur at fields of just a few teslas or even less, depending on the interlayer coupling strength. Moreover, the magnetic anisotropy in SAFs can be engineered through layer thickness, interface composition, or external stimuli (e.g., strain, electric field), providing additional tunability of the spin-flop behavior. This makes SAFs highly attractive for spintronic applications where field-tunable control over the Néel vector is desirable.

The qualitative behavior of spin-flop and spin-flip transitions remains the same in SAFs, but the tunable parameters and accessible energy scales make them more versatile platforms for fundamental research and applications alike.

2.2 Magnetic domains and domain walls

The concept of magnetic domains was introduced by Pierre-Ernest Weiss in 1906, marking the foundation of micromagnetics and domain theory¹¹⁶. It was developed to explain the behavior of ferromagnetic materials that exhibit zero net magnetization in the absence of an external field, yet reach saturation magnetiza-

tion ($\mu_0 M \approx 1$ T) under extremely weak applied fields (10^{-6} T), a response much stronger than in paramagnets. This behavior arises from the internal structure of FMs, which are divided into domains, regions where magnetic moments are aligned. Although each domain is magnetically ordered, the net magnetization may cancel out due to varying domain orientations. Application of a small external field causes reorientation of domains via DW motion, which is energetically more favorable than aligning individual spins.

Domain formation is primarily driven by minimization of the total magnetic free energy, which includes exchange, anisotropy, magnetostatic (demagnetizing), and Zeeman energy contributions^{7,8,III}. Figure 2.7 demonstrates the evolution of magnetic domain structure to minimize demagnetizing energy. To minimize the total magnetic free energy, the magnetic system adopts a multidomain configuration whose equilibrium structure is dictated by both intrinsic material properties (e.g., exchange stiffness, anisotropy) and extrinsic factors (e.g., sample shape, size, external field, and magnetic history)^{7,8,III}. By breaking into smaller domains and introducing closure domains (e.g., 90° rotations), magnetic surface and volume charges, and hence demagnetizing fields, can be minimized or eliminated. However, the presence of DWs introduces energy penalties due to violations of exchange and anisotropy energy preferences. The energy cost associated with creating DWs prevents the sample from breaking into too small domains^{7,8,III}.

The formation and structure of DWs arise from a balance between *exchange interaction* and *magnetic anisotropy*. In a FM, a DW separates domains, with spins gradually rotating from one domain orientation to another^{7,8,III}. The exchange energy associated with this rotation increases with the misalignment of neighboring spins. For a 180° domain wall spanning N atomic sites, the total exchange energy per unit area is approximately^{7,8,III}:

$$\sigma_E = \frac{JS^2\pi^2}{Na_0^2}, \quad (2.15)$$

where J is the exchange integral, S is the spin magnitude, and a_0 is the lattice constant. While exchange energy favors wide domain walls ($N \rightarrow \infty$), *magnetic anisotropy* counteracts this tendency. Rotating spins away from the easy axis incurs anisotropy energy, which prefers sharp transitions and thinner walls. For

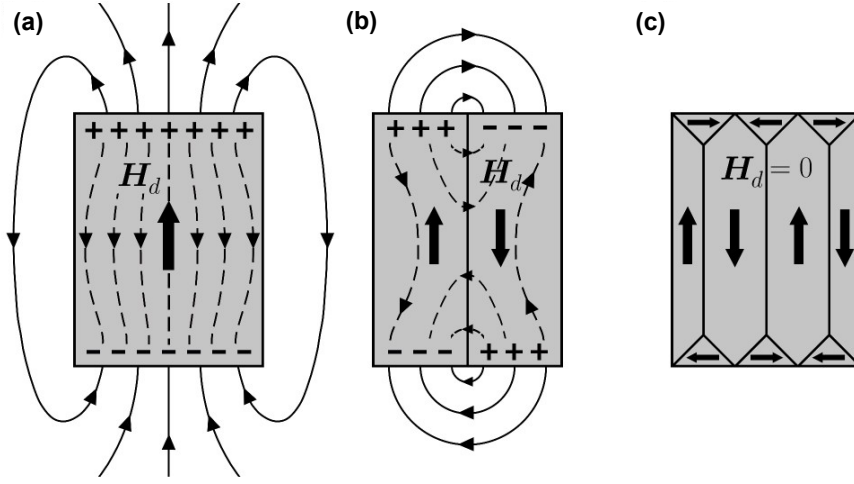


Figure 2.7: Evolution of magnetic domain structure to minimize demagnetizing energy. (a) A uniformly magnetized state produces strong demagnetizing fields due to surface magnetic charges. (b) Splitting into multiple domains reduces the net demagnetizing field by redistributing these charges. (c) The formation of closure domains with 90° rotated magnetization eliminates magnetic charges entirely, minimizing the total demagnetizing energy.

uniaxial anisotropy K_u , the anisotropy energy per unit area is:

$$\sigma_A = \frac{NK_u a_0}{2}. \quad (2.16)$$

The total domain wall energy density is then:

$$\sigma_{DW} = \frac{JS^2\pi^2}{Na_0^2} + \frac{NK_u a_0}{2}. \quad (2.17)$$

Minimizing σ_{DW} with respect to N yields the equilibrium domain wall width δ_{DW} and the corresponding energy:

$$\delta_{DW} = \pi \sqrt{\frac{A}{K_u}}, \quad \sigma_{DW} = \pi \sqrt{AK_u}, \quad (2.18)$$

where A is the exchange stiffness. There are two primary types of domain walls

(see figure 2.8)^{7,8,III}:

- **Bloch walls:** Magnetization rotates within the plane of the wall, typically stable in bulk due to minimal stray fields.
- **Néel walls:** Magnetization rotates perpendicular to the wall plane, favored in thin films to reduce demagnetizing energy.

In bulk FMs, *Bloch walls* are typically more stable because they produce no magnetic charges. For example, for a Bloch wall lying in the yz -plane, the magnetic charge density is given by $\rho_m = -\nabla \cdot \mathbf{M} = 0$. However, in thin films with in-plane magnetization, *Néel walls* are energetically preferred because the out-of-plane rotation of magnetization in Bloch walls would increase the demagnetizing energy. In reality, thin films are finite and contain imperfections (e.g., structural defects, grain boundaries), which act as nucleation centers for domain formation and play a crucial role in magnetization switching.

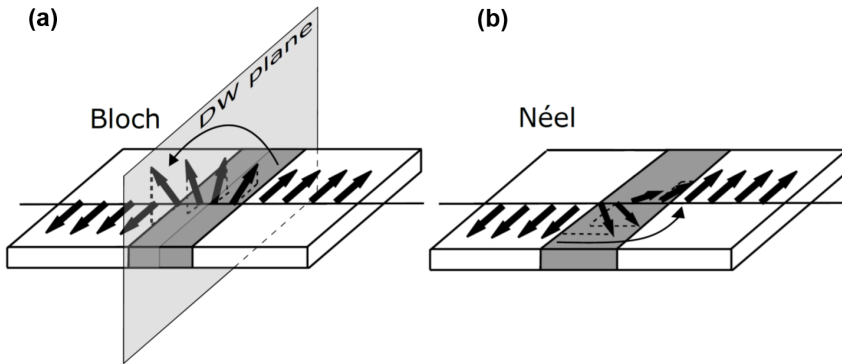


Figure 2.8: Internal structure of magnetic domain walls. (a) *Bloch wall*: the magnetization rotates within the plane of the domain wall. (b) *Néel wall*: the magnetization rotates in a plane perpendicular to the domain wall. Bloch walls are typically favored in bulk materials, while Néel walls are energetically preferred in thin films due to reduced demagnetizing energy (Reproduced from [7]).

2.2.1 Domains and domain walls in antiferromagnets and synthetic antiferromagnets

As previously discussed, the formation of magnetic domains in FMs is largely driven by the demagnetizing field. The additional energy cost of DWs is offset by the reduction in demagnetizing energy in the multidomain state. In contrast, the situation in AFs is fundamentally different. In zero external magnetic field, *compensated* AFs exhibit zero net magnetization and thus generate no demagnetizing fields. Nevertheless, experimental observations have confirmed the existence of domain structures in many AFs. In contrast to FMs, antiferromagnetic domains form naturally below the Néel temperature T_N as a consequence of spontaneous symmetry breaking, and symmetry-equivalent orientations of the Néel vector (e.g., $\pm \mathbf{L}$ and/or crystallographically equivalent easy-axis variants). Their formation is governed primarily by exchange and magnetic anisotropy energies, while the resulting domain patterns are strongly influenced by magnetoelastic coupling, structural defects, and sample boundaries rather than by magnetostatic energy.

The structure and energetics of DWs in AFs are governed by the same fundamental principles as in FMs. The rotation of the antiferromagnetically ordered spins across a domain boundary violates the exchange interaction, which favors a uniform spin alignment. This cost is counterbalanced by the magnetic anisotropy, which prefers spins aligned along specific directions. The equilibrium DW width and internal structure are thus determined by the competition between these two energy contributions. Unlike in FMs, where demagnetizing energy stabilizes domains, various other mechanisms are believed to be responsible for the existence and stability of AF domains. One proposed mechanism is *entropy stabilization*^{5,7}. In the absence of demagnetizing fields, DWs in AFs are not intrinsically strongly pinned and may thermally fluctuate around their equilibrium positions. Nevertheless, an antiferromagnetic DW can be pinned through interactions with another wall of the same helicity, or extrinsically by structural defects, disorder, or strain in real materials. This fluctuation increases the entropy S associated with the wall and lowers the free energy F . The change in free energy due to the presence of a DW is given by: $\Delta F = S_{DW}\sigma_{DW} - T\Delta S$, where S_{DW} is the domain wall area, σ_{DW} is the wall energy density, and T is the temperature. A multidomain state can thus become thermodynamically favorable ($\Delta F < 0$) if the domain wall energy is sufficiently low, which typically occurs in materials with weak magnetic

anisotropy, i.e., $\sigma_{\text{DW}} \sim \sqrt{K_u A}$.

Another key factor in AF domain formation is the role of *magneto-elastic interactions*⁵. At the Néel temperature (T_N), the transition into the antiferromagnetic phase generates a mismatch between bulk and surface magneto-elastic strains. This mismatch creates virtual "incompatibility elastic charges", which act analogously to magnetic charges in FMs and lead to an energy cost that can be minimized by forming domains. Finally, real AF materials contain lattice imperfections and defects that contribute to domain stabilization. These imperfections act as pinning centers for DWs, freezing in the domain structure established during the phase transition at T_N .

A key distinction between ferromagnetic and antiferromagnetic DW dynamics arises from the strong inter-sublattice exchange in AFs. In FMs, DW motion is accompanied by internal-angle precession, leading to Walker breakdown above a critical driving field and a corresponding reduction in DW velocity⁷. In contrast, in AFs, the strong exchange coupling rigidly locks the sublattices and suppresses internal-angle precession. As a result, Walker breakdown is absent, and AF DWs can in principle reach much higher velocities, ultimately limited by the magnon velocity⁷. Furthermore, uniform magnetic fields couple only weakly to compensated AFs, making field-driven DW motion inefficient. Instead, staggered fields or SOTs, which couple directly to the Néel vector, provide an efficient mechanism to drive AF DWs¹¹⁷.

2.3 Current-induced spin torques

Current-induced spin torques provide a powerful mechanism to manipulate and excite magnetic order in FMs, AFs, and SAFs¹¹⁸. Two main mechanisms are distinguished: STT, which arises from spin-polarized charge currents flowing through magnetic textures or multilayers^{104,118–123}, and SOT, which originates from SOC-mediated spin accumulation or spin currents and can emerge from effects such as the SHE or the interfacial iSGE (Edelstein effect)^{20,123–127}. In this thesis, SOT is the primary mechanism exploited for the manipulation of magnetic order in various SAF systems. In the following, a brief overview of STT and SOT is provided.

2.3.1 Spin-transfer torque

The concept of STTs revolutionized spintronics by enabling the control of magnetization using spin-polarized electric currents instead of external magnetic fields^{104,118–123}. This mechanism was first introduced by Berger and Slonczewski in the 1990s^{104,119,120}, who demonstrated that a flow of spin-polarized electrons can transfer angular momentum to a ferromagnet, exerting a torque on its magnetization.

At the microscopic level, STT originates from the exchange interaction between spin-polarized conduction electrons and the local magnetic moments in a FM, typically modeled by the *s-d exchange interaction*. A simplified physical picture considers the spin-split bands in a FM. If the spin polarization of the incoming electrons is *collinear* with the magnetization, both majority and minority spins are transmitted through the respective spin channels. Whereas, if the spin polarization is *non-collinear*, the transverse component of the spin must be absorbed by the magnetization to allow conduction. This absorption results in a torque on the local moments. This torque, known as the *Slonczewski torque* or *damping-like* (DL) torque, can be expressed as:

$$\tau_{\text{STT}} \propto \mathbf{M} \times (\mathbf{M} \times \mathbf{P}), \quad (2.19)$$

where $\mathbf{M}$ is the unit vector of the magnetization, and $\mathbf{P}$ is the spin polarization direction of the incoming current. This torque has the same form as the Gilbert damping term in the LLG equation and can either enhance or oppose damping depending on current polarity (detail discussion in [chapter III](#)).

When the STT compensates the natural damping, it can lead to either magnetization reversal or sustained precessional motion. These effects form the basis for various applications such as in STT based MRAM (STT-MRAM), a promising non-volatile memory technology offering fast switching, high endurance, and CMOS compatibility^{128,129} and in spin-torque nano-oscillators (STNOs) devices capable of generating microwave frequency signals through sustained magnetization precession^{73–78}. In practical devices, such as spin valves or MTJs, a typical STT structure consists of two ferromagnetic layers separated by a NM spacer ([figure 2.9](#)). The *reference layer* (FM1) is pinned (e.g., by exchange bias or increased thickness), while the *free layer* (FM2) is switchable via STT. For one current polar-

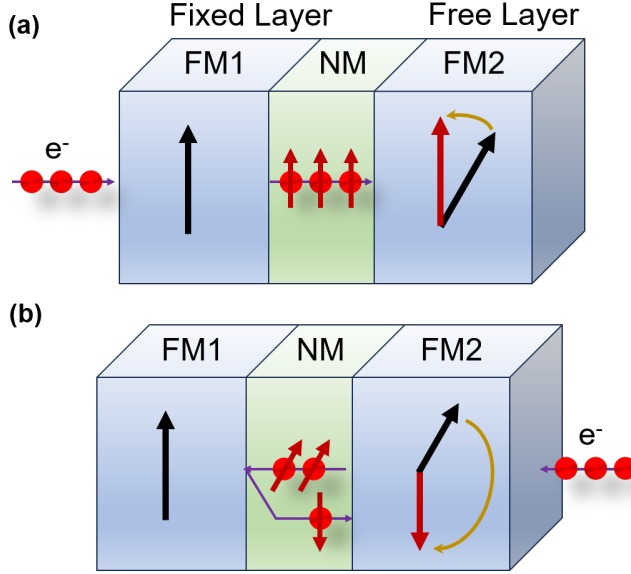


Figure 2.9: Schematic illustration of spin-transfer torque (STT) switching in a spin-valve structure. (a) A spin-polarized current flows from the reference (fixed) layer (FM1) to the free layer (FM2) via the NM spacer layer, with FM1 acting as a spin polarizer. The transferred spin angular momentum exerts a torque on FM2, aligning its magnetization parallel to that of FM1. (b) For the opposite current polarity, electrons flow from FM2 to FM1. Only spins aligned with FM1 are transmitted, while those with opposite spin are reflected. The reflected spins re-enter FM2, applying a torque that reorients its magnetization antiparallel to FM1.

ity, electrons flow from FM1 to FM2, exerting a torque that aligns FM2 *parallel* to FM1 (figure 2.9a). Similarly, for the opposite polarity, electrons flow from FM2 to FM1 and those electrons with spins antiparallel to FM1 are reflected, exerting a torque on FM2 that aligns it *antiparallel* to FM1 (figure 2.9b).

In addition to the even in magnetization DL torque of the form equation (2.19), experiments have revealed an additional component, known as the *field-like* (FL) torque, that is *odd* in magnetization and takes the form: $\tau_{\text{FL}} \propto \mathbf{M} \times \mathbf{P}$. This torque has the same symmetry as the precessional term in the LLG equation (more on chapter III). While the DL torque can be understood from angular momentum conservation, the origin of the FL torque is more complex and is still under investigation, with proposed mechanisms including Rashba SOC and interfacial

exchange fields. More generally, FL and DL (antidamping) torques can be viewed as steady-state processes arising from a continuous flow of spins into a FM, where the injected spins are misaligned with the exchange field. The resulting torque depends on the competition between the spin lifetime and the precession time around the exchange field: if the precession time is much longer than the spin lifetime, the steady-state spin density aligns with the injected spin polarization, yielding a FL torque; if the spin lifetime is much longer than the precession time, the steady-state spin polarization acquires a component perpendicular to both the magnetization and the incoming spin polarization, resulting in an DL torque.

2.3.2 Spin-orbit torque

While STT has been successfully implemented in technologies such as MTJs, its practical application is often limited by a fundamental trade-off: achieving efficient switching requires high current densities, which in turn necessitate low-resistance junctions to prevent device damage, yet reliable readout demands high magnetoresistance. Moreover, optimizing STT-based devices frequently involves intricate multilayer stacks, including over ten different layers, to fine-tune spin polarization, stabilize fixed magnetic layers, and minimize stray magnetic fields.

SOTs emerge as a compelling alternative, leveraging relativistic SOC in materials to generate torques on magnetic moments^{20,123–127}. These torques are induced by charge currents flowing through structures that lack inversion symmetry, where the interplay between SOC and exchange interaction allows angular momentum to be transferred from the lattice to the spin system. The spin-orbit interaction, already known to underpin phenomena such as magnetocrystalline anisotropy, AHE, anisotropic magnetoresistance (AMR), and spin relaxation, thus provides a versatile and efficient route for controlling magnetic dynamics. Importantly, SOTs enable a range of functionalities, including magnetization switching, DW and skyrmion motion, spin wave generation, and even operation in magnetic insulators and AFs¹²⁴. They also allow for decoupling of the read and write paths in devices, enhancing endurance and switching speeds compared to STT-based designs.

Two primary microscopic mechanisms underpin the generation of SOTs: the iSGE and the SHE. In bilayer systems, typically comprising a HM layer adjacent

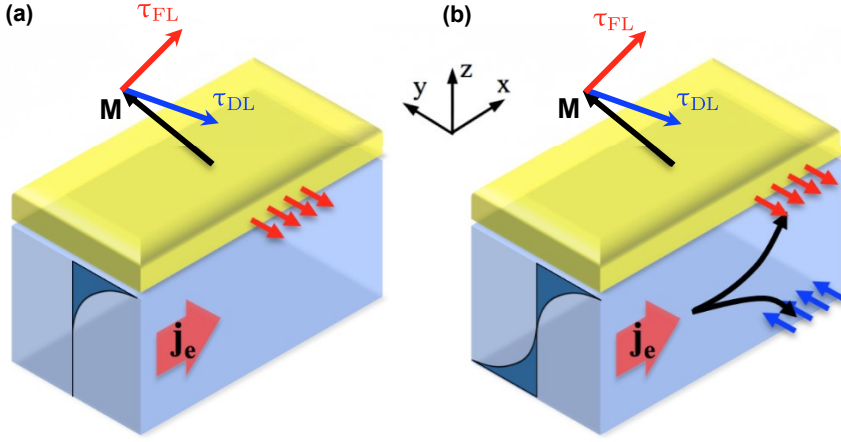


Figure 2.10: Two primary model mechanisms for spin-charge interconversion at the NM/FM interface: (a) the iSGE and (b) the SHE. Both mechanisms generate field-like (τ_{FL}) (red arrows) and damping-like (τ_{DL}) (blue arrows) torques on the FM. The small red and blue arrows indicate the nonequilibrium spin accumulation with polarization $\mathbf{P}$ at the interface, with their spatial distribution illustrated as shaded regions on the side of the structure. The large red and blue arrows represent the directions of the resulting torques. Adapted from [124]

to a FM, the origin of SOT is commonly attributed to a combination of these two effects. The iSGE-SOT arises from a non-equilibrium spin density generated locally at the FM/NM interface or within the FM itself due to SOC and broken inversion symmetry [see figure 2.10a]. This spin accumulation exerts a torque on the magnetization via exchange interaction. On the other hand, the SHE-SOT is attributed to a spin current generated in the bulk of the NM layer via the SHE, which then diffuses into the adjacent FM and exerts a torque [see figure 2.10b]. While distinct in origin, interface-driven vs. bulk-driven, these effects can coexist and contribute simultaneously to the observed torque. The following subsections outline the key principles behind each mechanism.

Inverse spin galvanic effect (Edelstein effect)

SOC arises when electrons move in an electric field, resulting in a Hamiltonian of the form:

$$\hat{\mathcal{H}}_{SO} \sim \hat{\mathbf{S}} \cdot (\hat{\mathbf{v}} \times \mathbf{E}), \quad (2.20)$$

where $\hat{\mathbf{S}}$ is the spin operator, $\hat{\mathbf{v}}$ is the velocity operator, and $\mathbf{E}$ is the electric field. This relativistic interaction is particularly strong in heavy elements such as platinum (Pt), where the large orbital velocities enhance SOC^{125,130,131}. Physically, SOC can be interpreted as the interaction of an electron's spin with the effective magnetic field that appears in its rest frame due to motion through an electric field. In atoms, where the Coulomb field is radially symmetric, the SOC is expressed as:

$$\hat{\mathcal{H}}_{\text{SO}} \sim \hat{\mathbf{S}} \cdot \hat{\mathbf{L}}, \quad (2.21)$$

where $\hat{\mathbf{L}}$ is the orbital angular momentum operator. In crystalline solids, where the Coulomb potential is periodic, the effective SOC becomes wavevector-dependent. As a result, the electronic energy bands depend on both spin and crystal momentum, i.e., $E^\sigma(\mathbf{k})$, with $\sigma = \uparrow, \downarrow$ ^{125,130,131}. The SOC Hamiltonian in crystals takes the form:

$$\hat{\mathcal{H}}_{\text{SO}} \sim \hat{\mathbf{S}} \cdot (\hat{\mathbf{k}} \times \mathbf{E}). \quad (2.22)$$

This Hamiltonian respects time-reversal symmetry, which implies:

$$E_{\text{SO}}^\uparrow(\mathbf{k}) = E_{\text{SO}}^\downarrow(-\mathbf{k}). \quad (2.23)$$

If spatial inversion symmetry is also present, then:

$$E_{\text{SO}}^\sigma(\mathbf{k}) = E_{\text{SO}}^\sigma(-\mathbf{k}), \quad (2.24)$$

which together imply $E_{\text{SO}}^\uparrow(\mathbf{k}) = E_{\text{SO}}^\downarrow(\mathbf{k})$, i.e., there is no SOC-induced spin splitting. Thus, *broken inversion symmetry* is essential for SOC to lift spin degeneracy. Inversion symmetry can be broken through various mechanisms, each giving rise to distinct spin-orbit effects. One common source is bulk inversion asymmetry (BIA), which arises from the intrinsic lack of a center of symmetry in certain crystal structures, such as the zinc-blende lattices found in many III-V and II-VI semiconductors¹³². Alternatively, structural inversion asymmetry (SIA) occurs in engineered heterostructures due to asymmetries in interface structure¹³³, doping gradients, or the presence of internal or external electric fields. In addition to these global symmetry-breaking mechanisms, local symmetry breaking caused by lattice defects or impurity atoms can also contribute to SOC effects. In low-dimensional systems, such as thin films, where electron motion is effectively confined to the plane, the interplay of BIA and SIA gives rise to two dominant forms of SOC: the Dresselhaus SOC¹³², which originates from BIA, and the Rashba SOC¹³³,

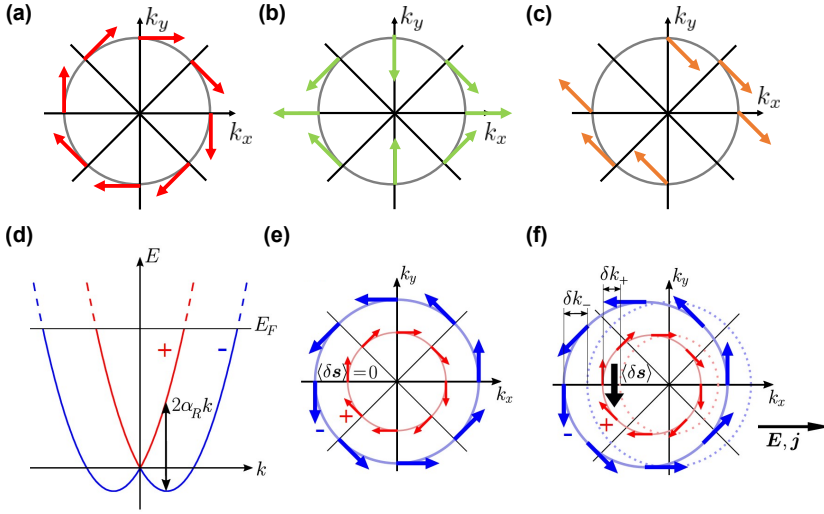


Figure 2.11: Effective spin-orbit fields and current-induced spin polarization. (a) Rashba spin-orbit field arising from structural inversion asymmetry (SIA). (b) Linear Dresselhaus spin-orbit field resulting from bulk inversion asymmetry (BIA). (c) Combined spin-orbit field due to the coexistence of Rashba and Dresselhaus interactions. (d) Spin-split conduction band structure induced by Rashba SOC. (e) Spin textures of Rashba-split bands at the Fermi level, showing symmetric spin polarization in equilibrium. (f) Non-equilibrium electron distribution under an applied electric field $\mathbf{E}$, leading to a net spin polarization $\langle \delta \mathbf{s} \rangle$ perpendicular to both the current density $\mathbf{j}$ and the SIA axis ($\hat{e}_z$). Adapted from [125].

which results from SIA. These two mechanisms lead to distinct spin-momentum locking configurations and effective spin-orbit fields that strongly influence spin dynamics and transport in the system. For example, in (001)-grown GaAs, the respective Hamiltonians are:

$$\hat{H}_R = \alpha_R(-k_x \sigma_y + k_y \sigma_x), \quad \hat{H}_D = \beta_D(k_x \sigma_x - k_y \sigma_y), \quad (2.25)$$

where α_R and β_D are the Rashba and Dresselhaus SOC strengths, and x, y, z correspond to the [100], [010], and [001] crystallographic directions. These SOC terms can be rewritten in a Zeeman-like form $\hat{\mathcal{H}}_{\text{SO}} = \mu_B \hat{\boldsymbol{\sigma}} \cdot \mathbf{B}_{\text{SO}}(\mathbf{k})$, resulting in effective spin-orbit fields:

$$\mathbf{B}_R(\mathbf{k}) = \frac{\alpha_R}{\mu_B} \begin{pmatrix} k_y \\ -k_x \end{pmatrix}, \quad \mathbf{B}_D(\mathbf{k}) = \frac{\beta_D}{\mu_B} \begin{pmatrix} k_x \\ -k_y \end{pmatrix}. \quad (2.26)$$

The textures of these effective spin-orbit fields around the Fermi surface (a circle in 2D) are illustrated in [figures 2.11a](#) and [2.11b](#). In equilibrium, the spin textures arising from the Rashba-split bands, with energies

$$E_{\pm}(\mathbf{k}) = \frac{\hbar^2 k^2}{2m^*} \pm \alpha_R |\mathbf{k}|, \quad (2.27)$$

cancel due to the symmetric occupation of states in $\mathbf{k}$ -space. However, when an external electric field is applied, this symmetry is broken, the Fermi surface is shifted in $\mathbf{k}$ -space, creating an asymmetric distribution of electrons¹²⁵. This results in a net non-equilibrium spin polarization, a phenomenon known as the *inverse spin galvanic effect* (iSGE), illustrated in [figures 2.11d](#) to [2.11f](#).

Quantitatively, the iSGE-induced spin density in a Rashba system under a current density $\mathbf{j}$ in the xy -plane is given by¹²⁵:

$$\langle \delta \mathbf{s} \rangle \approx \frac{m^* \alpha_R}{e \hbar E_F} \mathbf{j} \times \hat{e}_z, \quad (2.28)$$

where E_F is the Fermi energy and $\hat{e}_z$ is the unit vector perpendicular to the plane. The induced spin polarization is thus always orthogonal to both the direction of current flow and the axis of SIA.

Importantly, this non-equilibrium spin accumulation can exert a torque on the magnetization of an adjacent FM, AF or SAF layers. The mechanism involves the absorption of the transverse spin component at the interface, akin to STT. As discussed previously and illustrated in [figure 2.10a](#), the SOT generated via the iSGE can have both a FL and a DL component. Moreover, the iSGE can generate SOTs directly within a single FM layer, provided that the system possesses strong intrinsic SOC and lacks inversion symmetry¹²⁴. Similar effects have also been observed in AFs with appropriate crystal symmetry, extending the applicability of iSGE-induced torques beyond conventional FM/NM bilayer systems^{16,17,124}.

Spin Hall effect

In heterostructures consisting of a non-magnetic HM layer interfaced with a FM or AF layer, SOC can give rise to an additional mechanism that generates a pure spin current and a corresponding SOT: the *spin Hall effect* (SHE)^{20,134–137}. Analogous

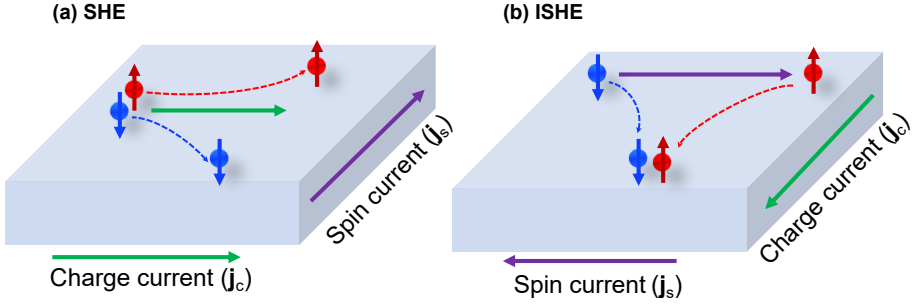


Figure 2.12: (a) Spin Hall effect (SHE) and (b) inverse spin Hall effect (ISHE). Solid green and purple arrows represent the direction of the charge current (j_c) and the spin current (j_s) respectively. The dotted arrows show the trajectories of spin-up and spin-down electrons. In the SHE, a longitudinal charge current generates a transverse spin current via spin-dependent deflection, while in the ISHE, a transverse spin current gives rise to a measurable charge voltage.

to the ordinary Hall effect (OHE), which deflects charge carriers perpendicular to both the electric field and an applied magnetic field, the SHE generates a spin current j_s transverse to the driving charge current j_c , even in the absence of an external magnetic field.

The direction of the spin angular momentum e_s carried by the spin current is orthogonal to both j_s and j_c , described by the relation²⁰

$$j_s = \alpha_{SH} e_s \times j_c, \quad (2.29)$$

where α_{SH} is the *spin Hall angle (SHA)*, which quantifies the efficiency and sign of the conversion from charge to spin current, and is determined by the strength of the spin-orbit interaction in the HM.

As illustrated schematically in [figure 2.12a](#), electrons with opposite spin orientations experience transverse deflections in opposite directions due to SOC, resulting in a pure spin current. Although the net spin density vanishes in the bulk due to symmetry, spin accumulation occurs at the lateral edges of the HM layer. This spin accumulation is partially counteracted by spin diffusion back into the bulk, leading to a steady-state non-equilibrium spin density near the HM surfaces¹³⁸. This process is also reversible, as shown in [figure 2.12b](#), via the *inverse spin Hall effect (ISHE)*, where a spin current can generate a transverse charge current^{20,139}.

Importantly, when the HM is in contact with an adjacent FM or AF layer, the spin current generated via the SHE can be injected across the interface and transferred to the magnetic moments in the adjacent layer, resulting in a SOT. This torque acts via the same fundamental spin-transfer processes as in the case of conventional STT, but without the need for a separate spin-polarizing FM layer. This is a key distinction: in SHE-SOT, the spin current originates intrinsically within the HM via SOC, rather than being sourced from a polarizing FM.

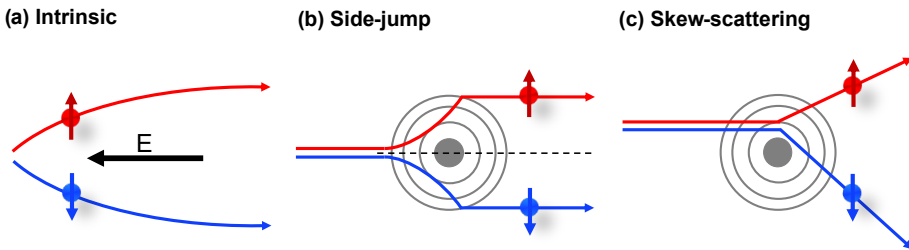


Figure 2.13: Mechanisms of SHE. (a) Intrinsic mechanism arising from the Berry curvature of the electronic band structure in systems with strong spin–orbit coupling. (b) Skew scattering, where spin-dependent deflection occurs due to asymmetric impurity scattering. (c) Side-jump scattering, involving a lateral displacement of electrons with opposite spins during scattering events, without a change in propagation direction.

The SHE has been observed experimentally through both extrinsic and intrinsic mechanisms²⁰. Figure 2.13 illustrates the three different possible mechanism responsible for the SHE. A landmark demonstration of the intrinsic spin Hall effect was provided by Wunderlich *et al.*¹³⁷, who used coplanar *p–n* diodes to detect circularly polarized electroluminescence at opposite edges of the spin Hall channel in two-dimensional hole gas (2DHG) devices. These experiments, conducted in the strong SOC regime, offered direct evidence for an intrinsic origin of the SHE. In comparison, Kato *et al.*¹⁴⁰ reported the extrinsic SHE in *n*-type GaAs using Kerr rotation microscopy to image spin accumulation caused by spin-dependent scattering. Together, these studies validated both intrinsic and extrinsic mechanisms and firmly established the SHE as a fundamental phenomenon in spintronics.

The extrinsic contributions to SHE arise from spin-dependent impurity scattering, as illustrated in figure 2.13b and figure 2.13c. In the *skew scattering* mechanism, electrons with opposite spin orientations are preferentially scattered in opposite directions due to the asymmetric spin-orbit interaction with charged impurities¹⁴¹. This can be viewed semi-classically: an electron moving in the electric

field $\mathbf{E}$ of an impurity experiences an effective magnetic field $\mathbf{B} \sim \mathbf{v} \times \mathbf{E}$ in its rest frame. This leads to spin-dependent deflection and results in a net transverse spin current. In contrast, the *side-jump* mechanism leads to a transverse displacement of the electron wave packet during scattering, without changing its propagation direction¹⁴². Again, electrons with opposite spins are shifted in opposite directions, giving rise to spin accumulation at the edges of the HM. After multiple scattering events, a significant spin imbalance builds up near the boundaries of the sample.

The intrinsic contribution to the SHE is fundamentally tied to the electronic band structure and its topological properties. In systems with strong intrinsic SOC, such as those with Rashba-type interactions, the electrons accelerated by an external electric field experience a time-dependent effective spin-orbit field. Due to the momentum dependence of this field, electrons acquire out-of-plane spin components while traversing the Brillouin zone^{20,143}. It is governed by the Berry curvature of the electronic bands and does not rely on disorder.

As discussed in Ref.^{20,124}, the SHE can give rise to both FL and DL torque components, τ_{FL} and τ_{DL} , respectively, much like in the case of iSGE-induced SOTs. Since both mechanisms can coexist in HM/FM or HM/AF bilayers, disentangling their contributions in experiments remains a challenge. Nevertheless, SHE-induced torques have been widely employed to switch magnetization in FM structures^{80,124} and, more recently, in antiferromagnets¹⁷, highlighting their technological relevance for next-generation spintronic devices.

2.4 Magnetoresistance effects

Magnetoresistance refers to the phenomenon where the electrical resistance of a material changes in response to an applied magnetic field or varies according to its magnetic ordering. This broad class of effects can be observed across a wide range of systems, including bulk non-magnetic materials, magnetically ordered compounds, and multilayered structures. Due to their significant potential in various technological applications, such as magnetic sensors and memory storage devices, magnetoresistance effects have been extensively investigated. This section aims to provide an overview of the principal magnetoresistive phenomena that arise from the interaction between electrical resistance and magnetic fields.

2.4.1 Anisotropic magnetoresistance

Anisotropic magnetoresistance (AMR) was first discovered in the 19th century by Lord Kelvin (William Thomson)¹⁴⁴. It is characterized by a relative change in the electrical resistance of a material depending on whether the magnetization vector is aligned parallel or perpendicular to the direction of the electric current^{144–146}. This phenomenon originates from spin-orbit coupling and is considered a bulk effect, typically observed in single-layer systems. AMR can be decomposed into two distinct contributions: crystalline and non-crystalline AMR¹⁴⁷. The *crystalline* contribution reflects the underlying symmetry of the crystal lattice, making it dependent on the angle between the magnetization direction and the crystallographic axes. This component is primarily observed in single-crystal samples, as it tends to average out in polycrystalline materials due to the random orientation of grains.

The *non-crystalline* contribution, historically observed by Lord Kelvin¹⁴⁴, is the more commonly encountered form of AMR. It describes the variation in resistance as a function of the angle between the magnetization and the current direction. The AMR ratio is typically expressed as:

$$\text{AMR} = \frac{R_{\parallel} - R_{\perp}}{R_{\text{av}}}, \quad (2.30)$$

where $R_{\parallel}$ and $R_{\perp}$ represent the resistances measured with the magnetization oriented parallel and perpendicular to the current, respectively, and R_{av} is their average value^{145,146}.

In addition to the longitudinal component, AMR also manifests as a transverse voltage signal, analogous to Hall effects. This transverse component is traditionally referred to as the *planar Hall effect (PHE)*. However, in the context of this thesis, it will be denoted as the *transverse AMR*, to emphasize that it arises from distinct symmetry considerations compared to conventional Hall phenomena.

The non-crystalline AMR is typically investigated using Hall bar devices through in-plane magnetic field rotations. Both longitudinal and transverse resistances exhibit characteristic angular dependencies given by:

$$R_{xx} = R_0 + \Delta R \cos 2\phi, \quad (2.31)$$

$$R_{xy} = \Delta R \sin 2\phi, \quad (2.32)$$

where ϕ is the angle between the magnetization vector and the direction of the electric current. An important property of AMR is that it is even with respect to magnetization, implying that the relevant quantity is the magnetization axis rather than its direction. As a result, AMR effects can persist even in systems with no net magnetization. In particular, AMR has been observed in collinear antiferromagnets^{148–150}, where magnetic sublattices align along the same axis but with opposing moments. In such cases, equations (2.31) and (2.32) remain valid, with ϕ defined as the angle between the Néel vector (the order parameter) and the current direction.

It should be noted that equations (2.31) and (2.32) represent the simplest angular dependence of the non-crystalline AMR. Higher-order harmonics may also be present, depending on the specific material and device geometry. For crystalline AMR, the angular dependence follows the symmetry of the crystal lattice.

2.4.2 Giant and tunneling magnetoresistance

The discovery of *GMR* in the late 1980s, independently reported by the groups of Albert Fert and Peter Grünberg, marked a pivotal moment in spintronics, ultimately earning them the Nobel Prize in Physics in 2007^{108–110}. GMR is observed in multilayer structures consisting of FM layers separated by a non-magnetic, conductive spacer. The term *giant* was coined due to the exceptionally large change in resistance, up to 50%, recorded in early experiments.

The origin of GMR lies in the dependence of electrical resistance on the relative alignment of the magnetizations of the ferromagnetic layers. It arises from spin-dependent electron scattering at the interfaces between magnetic and non-magnetic layers. This effect can be observed in two distinct configurations: current in plane (CIP) and current perpendicular to plane (CPP). The CPP geometry (the one used to describe STT in figure 2.9) is often considered more intuitive, as

the current flows perpendicularly through the layered structure. In this case, electrons become spin-polarized upon traversing the first FM layer. The non-magnetic spacer, typically chosen for its low SOC and long spin-diffusion length, allows the spin-polarized electrons to maintain their spin orientation with minimal scattering. When these electrons encounter the second FM layer, the degree of scattering depends on whether their spin is aligned with the magnetization of that layer. In an antiparallel configuration, enhanced scattering occurs, leading to higher resistance. Conversely, in a parallel configuration, scattering is reduced, and the resistance reaches a minimum.

In the CIP configuration, conduction electrons diffuse laterally across the stack. The magnetoresistive response in this case is governed by spin-dependent scattering events throughout the multilayer, and the phenomenon is typically analyzed using semiclassical models such as the Boltzmann transport equation, which emphasize interfacial spin scattering¹⁵¹.

A closely related phenomenon is *tunnel magnetoresistance (TMR)*, which emerges when the conductive spacer layer is replaced by a thin insulating barrier, and the current flows perpendicularly to the plane (CPP). In this regime, electron transport occurs via quantum tunneling, and the tunneling probability is influenced by the spin-dependent density of states in both ferromagnetic electrodes. As in GMR, the resistance depends on the relative orientation of the magnetizations: tunneling is more efficient in the parallel state than in the antiparallel one. Experimentally, TMR devices have demonstrated even larger resistance changes than their GMR counterparts, with reported values of up to 600%, compared to approximately 70% for GMR^{152,153}.

Both GMR and TMR are highly valued for technological applications due to their pronounced resistance modulation in response to magnetic configurations. A widely used implementation is the *spin-valve*, which serves as a fundamental building block in magnetic field sensors and non-volatile memory elements¹⁵⁴. A spin-valve typically consists of two ferromagnetic layers: a *reference layer*, whose magnetization is fixed, and a *free layer*, whose magnetization can be reoriented by external magnetic fields. The magnetic stability of the reference layer is usually ensured by selecting materials with differing coercivities or by employing an exchange bias mechanism.

2.4.3 Hall effects in magnetic systems

The Hall effect, first reported by Edwin Hall in 1879¹⁵⁵, describes the generation of a transverse voltage across a conductor subjected to a perpendicular magnetic field. This arises due to the Lorentz force acting on charge carriers, deflecting them orthogonally to both the current and magnetic field directions. The resulting charge imbalance at the sample edges gives rise to the well-known transverse voltage, known as the *ordinary Hall effect (OHE)*, which is present in all conducting materials.

Shortly after his initial discovery, Hall observed a significantly larger transverse voltage in ferromagnetic materials, which could not be explained solely by the Lorentz force. This enhanced response became known as the *anomalous Hall effect (AHE)*^{156,157}. Empirically, the total Hall resistivity in such materials is often expressed as:

$$\rho_{xy} = R_0 B_z + R_s \mu_0 M_z, \quad (2.33)$$

where R_0 is the ordinary Hall coefficient related to carrier concentration and B_z is the out-of-plane magnetic field. The second term represents the AHE, with R_s denoting the anomalous Hall coefficient and M_z the out-of-plane component of the magnetization.

While this expression provides a good empirical description for many systems, it does not account for AHE in materials with zero net magnetization, such as certain antiferromagnets. In particular, studies of non-collinear antiferromagnets^{33,34} and recent theoretical predictions for collinear systems such as RuO_2 ¹⁵⁸ have shown that a spontaneous Hall response can exist even in the absence of net magnetization. This challenges the traditional assumption that AHE is strictly proportional to M_z .

A microscopic understanding of the AHE was proposed by Karplus and Luttinger in 1954¹⁵⁹, who linked the effect to SOC through the concept of an *anomalous velocity* arising from relativistic corrections to the electron's Hamiltonian in a periodic potential. This anomalous velocity was later found to be directly related to the Berry curvature in momentum space^{157,160}. The Berry curvature, acting as a fictitious magnetic field in reciprocal space, can become strongly enhanced

near band crossings or degeneracies¹⁵⁷, offering a pathway to engineer the Hall response by tuning the Fermi level near such features.

Symmetry plays a critical role in determining whether a material can host a finite Berry curvature. Under time-reversal ($\mathcal{T}$) and inversion ($\mathcal{P}$) symmetries, the Berry curvature $\mathbf{\Omega}(\mathbf{k})$ transforms as $\mathcal{T}: \mathbf{\Omega}(\mathbf{k}) \rightarrow -\mathbf{\Omega}(-\mathbf{k})$ and $\mathcal{P}: \mathbf{\Omega}(\mathbf{k}) \rightarrow \mathbf{\Omega}(-\mathbf{k})$. Thus, in systems with combined $\mathcal{PT}$ symmetry, the Berry curvature must vanish¹⁶¹. This implies that materials such as the $\mathcal{PT}$ -symmetric antiferromagnet CuMnAs are not expected to exhibit an anomalous Hall response.

In the bulk, in addition to the intrinsic Berry curvature mechanism, two extrinsic scattering-based mechanisms, *skew scattering* and *side jump*, also contribute to the AHE^{141,142,157}. These effects are associated with asymmetric scattering of electrons at impurities due to SOC and become significant depending on the level of disorder and impurity concentration in the system.

Another related phenomenon is the *topological Hall effect (THE)*, which arises from real-space noncoplanar spin textures¹⁶². In this case, conduction electrons pick up a Berry phase as they move adiabatically through a topologically nontrivial magnetic structure, such as a skyrmion lattice¹⁶². The THE was first observed in MnSi, where it was associated with a skyrmion phase stabilized under certain field and temperature conditions¹⁶³. More complex magnetic configurations, such as those found in Mn₅Si₃ and the real space non-coplanar spin texture at low temperature in Mn₃Sn^{164,165}, have also been shown to produce a topological Hall signal. The underlying mechanism is governed by the scalar spin chirality:

$$\kappa = \vec{\mathbf{n}}_i \cdot (\vec{\mathbf{n}}_j \times \vec{\mathbf{n}}_k), \quad (2.34)$$

where $\vec{\mathbf{n}}_i$, $\vec{\mathbf{n}}_j$, and $\vec{\mathbf{n}}_k$ are unit vectors representing local spin directions at three neighboring sites within a triangular plaquette. Coplanar spin configurations yield $\kappa = 0$, while noncoplanar ones produce a finite chirality, giving rise to the THE.

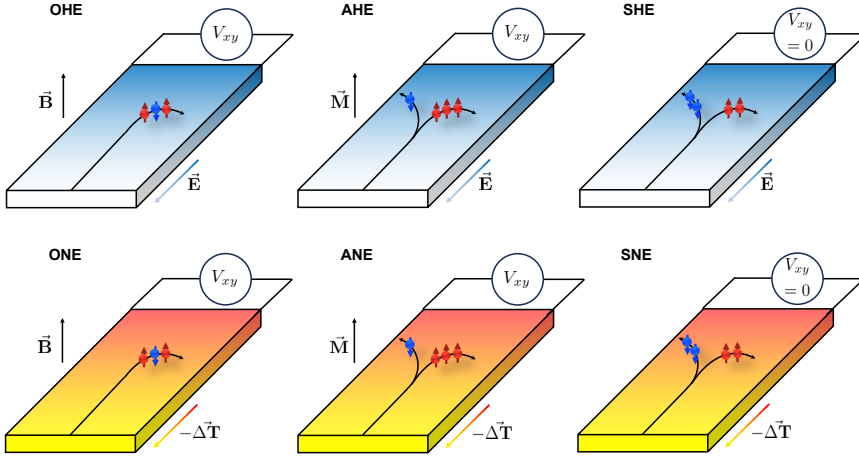


Figure 2.14: Schematics of the Hall and Nernst effects.

2.5 Thermoelectric effects

While the previously discussed magnetotransport phenomena are driven by an applied electric field or a potential gradient, similar effects can also arise when charge carriers are driven by a temperature gradient. These thermally induced phenomena often exhibit symmetries analogous to their electrically driven counterparts and are therefore commonly referred to as their thermal analogues. Importantly, a thermal gradient can couple not only to the charge current but also to the spin current, giving rise to a variety of thermoelectric and spin-dependent effects. The interplay between heat, charge, and spin transport forms the basis of the field known as *spin caloritronics*¹⁶⁶.

Spin caloritronics encompasses phenomena such as the spin-dependent Seebeck and Peltier effects, as well as thermal counterparts of Hall-type effects, collectively referred to as Nernst effects¹⁶⁶. Although thermoelectric signals are often treated as parasitic artefacts in magnetotransport experiments, they can also serve as sensitive probes of magnetic order. In this section, we provide a brief overview of some key thermoelectric effects relevant to magnetic and spintronic systems.

The most fundamental thermoelectric phenomenon is the Seebeck effect (or

thermopower), which describes the build-up of an electric voltage in response to a temperature gradient across a material. The Seebeck coefficient, S , characterizes this response and is defined as:

$$S = -\frac{\Delta V}{\Delta T}, \quad (2.35)$$

where ΔV is the voltage generated by a temperature difference ΔT .

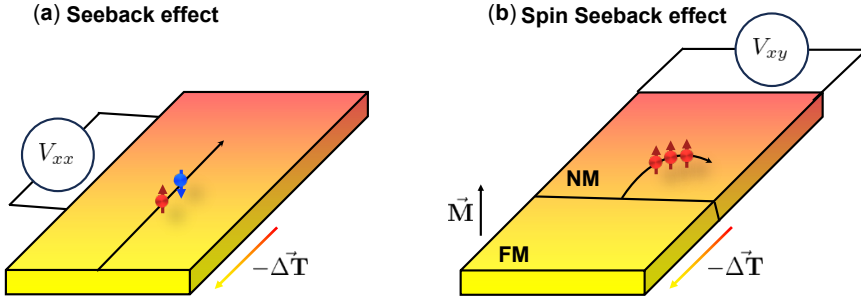


Figure 2.15: Schematic illustrations of the Seebeck effect and the spin Seebeck effect (SSE). The SSE shown corresponds to the longitudinal configuration, where the spin current and the thermal gradient are aligned. A transverse configuration also exists but is not depicted here.

Thermally induced charge currents can also interact with external magnetic fields and magnetic order, giving rise to the so-called Nernst effects. Analogous to the ordinary and anomalous Hall effects (OHE and AHE), the thermal counterparts are known as the ordinary Nernst effect (ONE) and the ANE, respectively. The ONE arises from the influence of an external magnetic field on a thermally driven charge current, while the ANE stems from the intrinsic magnetic structure of the material. In FMs and certain AFs, the ANE is understood to originate from the Berry curvature in momentum space³⁵, similarly to the AHE.

An analogue of the AMR can also be observed in the thermoelectric regime. In magnetic systems, the magnitude of the thermally induced charge current depends on the relative orientation between the magnetic order and the current direction. This effect, often referred to as anisotropic magneto-thermopower, shares the same symmetry as conventional AMR and can be detected in both longitudinal and transverse measurement geometries.

Beyond charge-based thermoelectric effects, spin-dependent thermal transport phenomena are also observed. For instance, the spin Nernst effect (SNE) has been reported in non-magnetic materials^{167,168}. It is the thermal analogue of the SHE and describes the generation of a transverse spin current perpendicular to a thermal gradient. The SNE provides a mechanism for generating pure spin currents via thermal means in materials with strong spin-orbit coupling.

A particularly notable phenomenon in this context is the spin Seebeck effect (SSE)^{169,170}. The SSE refers to the generation of a spin current in a ferromagnetic material due to a thermal gradient. This spin current is subsequently converted into a measurable charge current in an adjacent non-magnetic layer via the ISHE. Interestingly, the SSE cannot be fully explained by spin-polarized conduction electrons alone, as it has also been observed in insulating ferromagnets such as yttrium iron garnet (YIG). In such cases, the effect is attributed to the transport of magnons, quanta of spin-wave excitations, highlighting the role of collective spin excitations in spin caloritronics.

Figure 2.14 schematically compares the Hall and Nernst effects, illustrating the analogous nature of electrical and thermal responses in magnetic systems. Figure 2.15 illustrates the longitudinal geometry of the SSE, in which the spin current is nominally aligned with the temperature gradient. However, this configuration has never been experimentally verified, as it would require a spin diffusion length comparable to the sample thickness. A transverse geometry, the only one which has been verified, also exists, where the spin current flows perpendicular to the thermal gradient. In conducting FMs, this transverse configuration enables the detection of SSE signals with minimal interference from other thermoelectric effects such as the ANE.



Spin dynamics in antiferromagnets

This chapter develops the theoretical foundations necessary to understand the dynamic properties of AFs. We begin in [section 3.1](#) with an introduction to the LLG equation, which provides the fundamental framework for describing the time evolution of magnetic moments. This is followed in [section 3.2](#) by a concise overview of FMR, highlighting the key physical principles that serve as a reference point for more complex antiferromagnetic dynamics. In [section 3.3](#), we present a brief introduction to antiferromagnetic resonance (AFMR), accompanied by a detailed analysis of the characteristic eigenmodes that arise in both easy-axis and easy-plane anisotropy AFs. Subsequently, [section 3.4](#) offers a short discussion of spin waves (magnons). Finally, [section 3.5](#) introduces current-induced, SOT-driven self-oscillations in both FMs and AFs, summarizing the state-of-the-art developments. Collectively, these sections establish a coherent theoretical basis for the dynamic magnetic phenomena explored in the experimental chapters.

Keywords: antiferromagnetic dynamics, antiferromagnetic resonance (AFMR), magnons, chiral eigenmodes, SOT-driven nano-oscillators.

3.1 Magnetization dynamics and LLG equation

The dynamics of magnetization in ferromagnetic materials can be described by the interaction between the magnetization vector $\mathbf{M}$ and an effective magnetic field $\mu_0 \mathbf{H}_{\text{eff}}$. The effective magnetic field may include contributions from external magnetic fields, magnetic anisotropy, exchange interactions, and demagnetizing fields as described in the [section 2.1](#). When subjected to such an effective field, the magnetization experiences a torque given by the vector cross product $\mathbf{M} \times \mu_0 \mathbf{H}_{\text{eff}}$ ^{8,115}. This torque results in a precessional motion of the magnetization around the direction of the effective field. The time evolution of $\mathbf{M}$ under the influence of this field torque can be obtained by the Landau-Lifshitz equation^{115,171,172}:

$$\frac{d\mathbf{M}}{dt} = -\gamma \mathbf{M} \times \mu_0 \mathbf{H}_{\text{eff}}, \quad (3.1)$$

where μ_0 is the vacuum permeability and γ is the gyromagnetic ratio defined as

$$\gamma = \frac{g\mu_B}{\hbar}, \quad (3.2)$$

with g being the Landé g -factor, μ_B the Bohr magneton, and $\hbar = h/2\pi$ the reduced Planck's constant. The quantity γ sets the proportionality between the magnetic moment and its angular momentum. Thus, in the absence of any dissipative mechanism, the magnetization follows a circular trajectory around $\mu_0 \mathbf{H}_{\text{eff}}$, with no change in the magnitude of $\mathbf{M}$ or its projection onto $\mu_0 \mathbf{H}_{\text{eff}}$. A schematic illustration of the precessional motion of the magnetization vector is shown in [figure 3.1a](#). The angular frequency ω of this precessional motion is determined by the strength of the effective magnetic field and is given by^{8,115}

$$\omega = \gamma \mu_0 |\mathbf{H}_{\text{eff}}|. \quad (3.3)$$

From experimental observation, it is well established that the magnetization vector $\mathbf{M}$ does not precess indefinitely around the magnetic field, but rather tends to spiral and eventually align with it over time. For example, when a sufficiently strong magnetic field is applied to a ferromagnetic material, it can reach magnetic saturation i.e., a state in which all magnetic domains are aligned parallel to the external field direction. This behavior indicates that [equation \(3.1\)](#), which describes purely conservative precession, is incomplete; it lacks a mechanism that

drives the magnetization into alignment with the effective magnetic field ($\mu_0 \mathbf{H}_{\text{eff}}$). To account for this behavior, T. Gilbert introduced in 1955 a phenomenological damping term, which leads to the LLG equation ^{*174}:

$$\frac{d\mathbf{M}}{dt} = -\gamma \mathbf{M} \times \mu_0 \mathbf{H}_{\text{eff}} + \frac{\alpha}{M_{\text{sat}}} \left(\mathbf{M} \times \frac{d\mathbf{M}}{dt} \right), \quad (3.4)$$

where: α is the dimensionless Gilbert damping parameter and M_{sat} is the saturation magnetization. The first term on the right-hand side of Eq. (3.4) represents the torque responsible for the precession of the magnetization vector, while the second term introduces damping, causing the magnetization to spiral inward and eventually align with the effective field direction ($\mathbf{M} \parallel \mu_0 \mathbf{H}_{\text{eff}}$). A schematic illustration of the damped precessional motion of the magnetization vector is shown in figure 3.1b, where the trajectory of $\mathbf{M}$ converges toward the equilibrium direction defined by $\mathbf{H}_{\text{eff}}$.

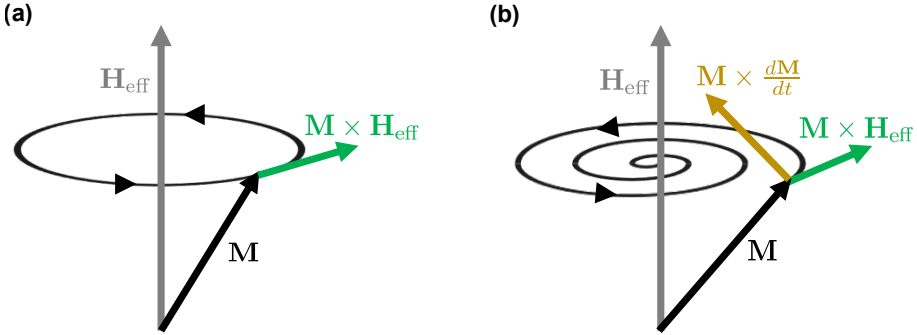


Figure 3.1: (a) Precessional motion of the magnetization $\mathbf{M}$ (black arrow) in an effective magnetic field $\mathbf{H}_{\text{eff}}$ as described by equation (3.1). The torque originating from the externally applied field is indicated in green. (b) Damped precessional motion of $\mathbf{M}$ under the action of the Gilbert damping torque (equation (3.4)). The damping torque, which leads to an alignment of the magnetization to the equilibrium position ($\mathbf{M} \parallel \mathbf{H}_{\text{eff}}$), is shown in golden arrow.

The damping parameter α encompasses all processes through which the system loses energy and relaxes toward equilibrium. It can be broadly categorized

^{*}Apart from the LLG equation, another equation that describes magnetization at finite temperature called the Landau-Lifshitz-Bloch (LLB) exists¹⁷³. In LLB equation, two different kinds of relaxation mechanisms are proposed where the magnitude of $\mathbf{M}$ is not conserved. In this thesis we only consider the LLG equation only.

into intrinsic and extrinsic contributions¹¹⁵. Intrinsic damping mechanisms are material-specific and include, for instance, SOC and energy transfer to the crystal lattice via phonon interactions. Extrinsic damping, on the other hand, arises from sample-dependent factors such as interface roughness, structural defects, impurities, and grain boundaries. Equation (3.4) forms the cornerstone of micromagnetic theory and provides a quantitative framework for describing a broad spectrum of magnetization dynamics phenomena, including ferromagnetic resonance, SW propagation, DW motion, and ultrafast switching dynamics in nanoscale systems.

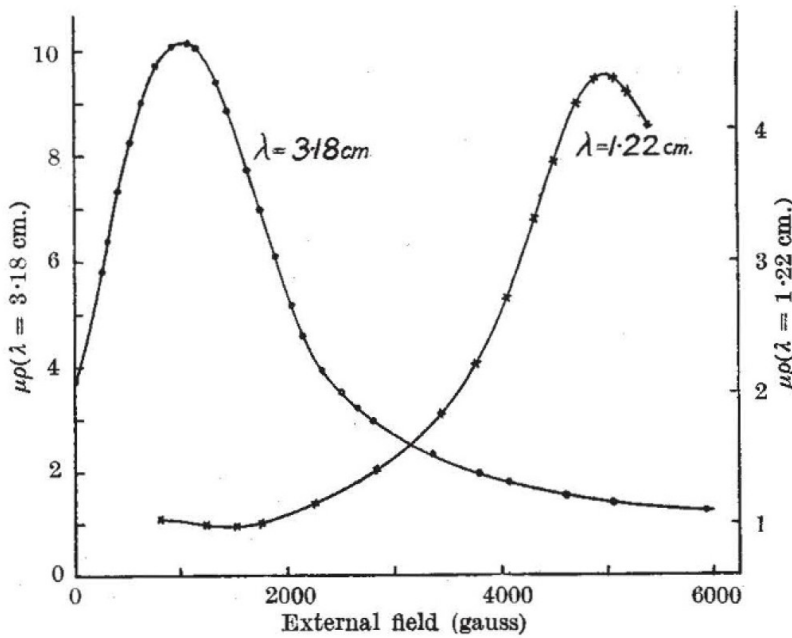


Figure 3.2: Ferromagnetic resonance (FMR) peak of a ferromagnetic metal (Ni) film. The vertical axis represents the product of the electrical resistivity ρ and the magnetic permeability μ of the conductor, which characterizes the energy loss of the microwave. The horizontal axis corresponds to the applied external magnetic field. The symbol λ denotes the wavelength of the microwave. Reproduced from¹⁷⁵.

3.2 Ferromagnetic resonance

FMR is a powerful spectroscopic technique used to investigate the dynamic response of magnetization in ferromagnetic materials under the influence of external magnetic fields. First observed experimentally by Griffiths in 1946¹⁷⁵ and later confirmed by Yager et al.¹⁷⁶, the theoretical framework explaining these observations was formulated by Kittel in 1947^{177,178}. Figure 3.2 shows the FMR spectrum of a Ni film as reported by Griffith¹⁷⁵. Below some key Kittel resonance conditions are highlighted for a thin ferromagnetic film, while more comprehensive treatments are available in earlier works^{115,179,180}.

FMR involves the resonant absorption of microwave radiation by a ferromagnetic material when the frequency of the applied microwave field matches the natural precession frequency of the magnetization $\mathbf{M}$ (see figure 3.3a)^{115,172}. In the presence of a static field $\mathbf{H}_0$ applied along the z -axis and a transverse radio frequency (rf) field $\mathbf{h}_{\text{rf}}(t) = \mathbf{h}_0 e^{i\omega t}$, the dynamics of $\mathbf{M}$ are governed by the LLG equation of the form equation (3.4), where the effective field $\mathbf{H}_{\text{eff}} = \mathbf{H}_0 + \mathbf{H}_A - \mathcal{N}\mathbf{M}$ includes external, anisotropy, and demagnetizing contributions. When the frequency of $\mathbf{h}_{\text{rf}}$ matches the natural precessional frequency of $\mathbf{M}$, resonance occurs, allowing efficient transfer of energy from the microwave field into the magnetic system. At resonance, microwave absorption increases sharply, detected by sweeping the static magnetic field at a fixed microwave frequency. Maximum absorption occurs when the magnetization precession is phase-aligned with the driving field. As the field moves away from resonance, phase mismatch reduces absorption, producing a characteristic spectrum whose peak and linewidth reveal damping and material properties.

Assuming small-angle precession such that the time-dependent part is smaller than the time-independent one ($\mathbf{m}(t) \ll \mathbf{M}_0$) and neglecting higher-order terms, the linearized response of the transverse magnetization $\mathbf{m} = (m_x, m_y)$ to the driving field $\mathbf{h}_{\text{rf}}$ is expressed via the Polder susceptibility tensor $\bar{\chi}$ ^{115,181}:

$$\mathbf{m} = \bar{\chi} \mathbf{h}_{\text{rf}}, \quad \text{with } \bar{\chi} = \begin{pmatrix} \chi_{xx} & i\chi_{xy} \\ -i\chi_{xy} & \chi_{yy} \end{pmatrix}. \quad (3.5)$$

The second rank Polder susceptibility tensor $\bar{\chi}$ quantifies the linear response of

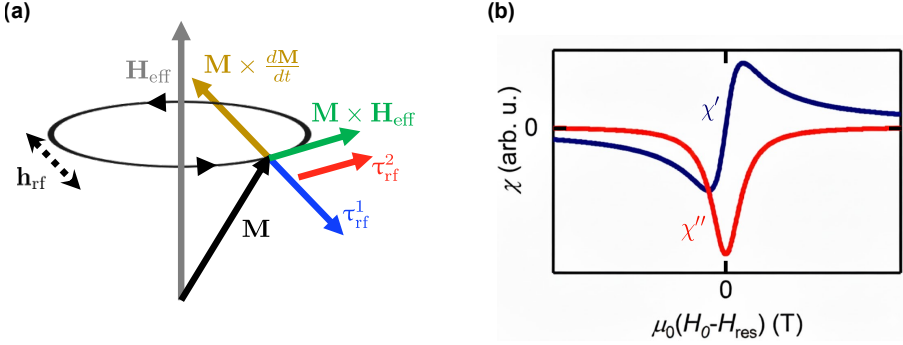


Figure 3.3: (a) Precessional motion of the magnetization $\mathbf{M}$ in an effective magnetic field $\mathbf{H}_{\text{eff}}$, as described by [equation \(3.4\)](#). The torque originating from the externally applied field is shown in green, while the two torques from the rf field are denoted as τ_{rf}^1 and τ_{rf}^2 in blue and red, respectively. (b) Real (blue) and imaginary (red) parts of the high-frequency susceptibility χ as a function of the externally applied magnetic field $\mu_0 H_0$.

the magnetization to a small rf magnetic field perpendicular to its equilibrium direction. The susceptibility tensor elements can be decomposed into real and imaginary parts: $\chi = \chi' + i\chi''$, where the real part χ' describes the dispersion of the system, and the imaginary part χ'' accounts for energy absorption¹¹⁵. The frequency dependence of χ'' exhibits a Lorentzian lineshape, with the maximum corresponding to the resonance condition, where energy absorption is most efficient. The typical lineshape of the susceptibility is shown in [figure 3.3b](#).

To determine the resonance condition for FMR, one evaluates the condition $\det(\tilde{\chi}^{-1}) \stackrel{!}{=} 0$, which leads to the Kittel equation for the resonance frequency^{177,178}. For thin film geometries, where the thickness is typically on the nanometer scale, simplified forms of the Kittel equation can be obtained for specific field orientations. When the static magnetic field is applied in the plane of the film (in-plane geometry), Kittel resonance condition is given by:

$$f = \frac{\gamma}{2\pi} \mu_0 \sqrt{(H_0 + H_A)(H_0 + H_A + M_{\text{sat}})}. \quad (3.6)$$

Conversely, in the out-of-plane geometry, where $\mathbf{H}_0$ is oriented perpendicular to the film plane, the anisotropy field and saturation magnetization are combined into an effective magnetization $M_{\text{eff}} = M_{\text{sat}} - H_A$, since these two quantities cannot be independently resolved in standard FMR measurements. The resonance

condition becomes:

$$f = \frac{\gamma}{2\pi} \mu_0 (H_0 - M_{\text{eff}}). \quad (3.7)$$

The diagonal elements of the susceptibility tensor also contain an imaginary component proportional to the Gilbert damping parameter α , which reflects the dissipative processes in the magnetization dynamics. Experimentally α is extracted from the full-width-at-half-maximum (FWHM) linewidth ($\mu_0 \Delta H_{\text{FWHM}}$). The total linewidth is a combination of intrinsic damping and extrinsic inhomogeneous broadening (ΔH_{inh}) and the full expression is given by:

$$\mu_0 \Delta H_{\text{FWHM}} = \mu_0 \Delta H_{\text{inh}} + \frac{4\pi\alpha f}{\gamma}. \quad (3.8)$$

The physical origin of damping mechanisms can vary and often depends on the specific material system. Common contributions include two-magnon scattering, eddy currents in conducting films, and spin-pumping effects at interfaces. The frequency-independent inhomogeneous linewidth broadening arises from spatial variations in the magnetic properties across the sample. As such, the linewidth serves as an important diagnostic of both intrinsic and extrinsic damping phenomena in ferromagnetic materials.

3.3 Antiferromagnetic resonance

The phenomenon of AFMR was first explored theoretically in the pioneering works of Kittel^{182,183} and Nagamiya¹⁸⁴, and it was subsequently confirmed experimentally by Ubbink et al. in 1952¹⁸⁵. Since this initial discovery, extensive research has been conducted to understand AFMR across a wide range of materials exhibiting diverse crystallographic and magnetic structures^{5,96,186–188}. In this section, we extend the theoretical framework previously developed for FMR (section 3.2) to the case of AFMR. Our goal is to highlight both the characteristic features of AFMR and the subtle distinctions that set it apart from its ferromagnetic counterpart, particularly in terms of the eigenmodes observed in AF systems. For clarity and easy understanding, we focus on the simplest case: a two-sublattice AF with equivalent magnetic sublattices, considering two typical anisotropy scenarios: easy-axis and easy-plane. My thesis explores both scenarios within SAF heterostructures. Spe-

cifically, I examine SOT-driven DW propagation in SAFs with uniaxial magnetic anisotropy oriented perpendicular to the film plane. Additionally, I investigate AFMR induced by SOT, as well as auto-oscillations in antiferromagnetic oscillators realized in systems with easy-plane magnetic anisotropy. This simplified two-sublattice model also provides valuable insight into the eigenmode structure of SAFs, the focus of the present thesis, which can be effectively regarded as systems composed of a coupled antiferromagnetic sublattices. We directly examine the nature and type of eigenmodes excited in AFMR, along with their resonance frequency dependence on the magnitude and orientation of an applied magnetic field. For a more rigorous and detailed theoretical treatment, the reader is referred to the relevant literatures^{115,187}.

The theoretical treatment of AFMR is inherently more intricate than that of ferromagnetic resonance, as it involves the coupled dynamics of the two sublattice magnetizations, $\mathbf{M}_1$ and $\mathbf{M}_2$, each governed by its own equation of motion and strongly coupled via the exchange interaction. To describe these dynamics, we employ the LLG equation for each magnetic sublattice (or ferromagnetic layer in the case of SAFs), where the unit magnetization vectors $\mathbf{M}_i$ evolve under the influence of effective interaction fields $\mathbf{H}_{\text{eff},i}$.

$$\frac{d\mathbf{M}_i}{dt} = -\gamma \mathbf{M}_i \times \mu_0 \mathbf{H}_{\text{eff},i} + \frac{\alpha}{M_{\text{sat}}} \left(\mathbf{M}_i \times \frac{d\mathbf{M}_i}{dt} \right), \quad (3.9)$$

where, we have considered for the two equivalent sublattices, $\gamma_1 = \gamma_2 = \gamma$, $M_{\text{sat},1} = M_{\text{sat},2} = M_{\text{sat}}$, $\alpha_1 = \alpha_2 = \alpha$, and the effective field is defined as the sum of the applied field ($\mathbf{H}_0$), anisotropy field ($\mathbf{H}_A$), and the interlayer exchange field ($\mathbf{H}_E$) as:

$$\mathbf{H}_{\text{eff},i} = \mathbf{H}_0 + \mathbf{H}_A + H_E \mathbf{M}_j \quad (3.10)$$

where $i, j = 1, 2$ correspond to the two magnetic sublattices.

3.3.1 Eigenmodes in easy-axis antiferromagnet

Now, let us consider the case of an easy-axis AF with uniaxial anisotropy oriented along the z -axis. Under this condition, the coupled LLG equations given in

equation (3.9) take the form¹⁸³:

$$\frac{d\mathbf{M}_1}{dt} = -\gamma \mathbf{M}_1 \times \mu_0 (\mathbf{H}_0 + \mathbf{H}_A + H_E \mathbf{M}_2) + \frac{\alpha}{M_{\text{sat}}} \left(\mathbf{M}_1 \times \frac{d\mathbf{M}_1}{dt} \right), \quad (3.11)$$

$$\frac{d\mathbf{M}_2}{dt} = -\gamma \mathbf{M}_2 \times \mu_0 (\mathbf{H}_0 - \mathbf{H}_A + H_E \mathbf{M}_1) + \frac{\alpha}{M_{\text{sat}}} \left(\mathbf{M}_2 \times \frac{d\mathbf{M}_2}{dt} \right). \quad (3.12)$$

The sign difference in the uniaxial anisotropy term reflects the antiparallel alignment of the two sublattice magnetizations, with the anisotropy field pointing along $+z$ for $\mathbf{M}_1$ and along $-z$ for $\mathbf{M}_2$ ¹⁸³. To explore the dynamical response of AFMR in the presence of a small, time-varying perturbing rf magnetic field $\mathbf{h}_{\text{rf}}(t)$, we follow the same linearization procedure as demonstrated for FMR in section 3.2. Specifically, the magnetization vectors and effective fields are expressed as the sum of steady (dc) and oscillatory (ac) components. These expressions are then substituted into the coupled LLG equations in equations (3.11) and (3.12). Assuming that the ac components are small compared to the steady ones ($\mathbf{m}_i(t) \ll \mathbf{M}_{0,i}$), we linearize the system and solve for the eigenfrequencies of the resulting dynamics. The analytical expressions for the resonance frequencies are summarized in the table 3.1 for two representative configurations: one where the applied magnetic field is aligned parallel to the easy axis, and another where it is applied perpendicular to the easy axis.

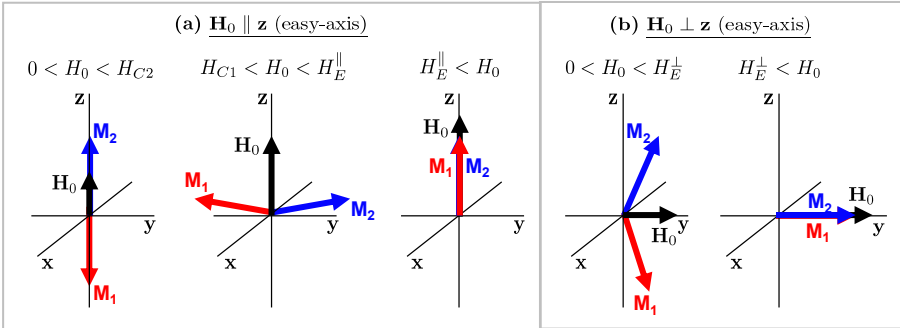


Figure 3.4: Ground-state configurations of a two-sublattice AF with an easy-axis anisotropy along the z-direction for (a) a magnetic field applied parallel to the easy axis ($\mathbf{H}_0 \parallel \mathbf{z}$) and (b) a magnetic field applied perpendicular to the easy axis ($\mathbf{H}_0 \perp \mathbf{z}$). Adapted from [115].

The ground states of a two-sublattice AF with easy-axis anisotropy are shown in figure 3.4 for a magnetic field ($\mathbf{H}_0$) applied either parallel or perpendicular

Table 3.1: Eigenfrequencies of a two-sublattice AF with easy-axis anisotropy ($H_A > 0$) at $T = 0K$. Reproduced from [115].

Field direction ^a	Ground state	Field range ^b	Eigenfrequencies ^c
$\mathbf{H}_0 \parallel \mathbf{z}$	Antiparallel	$0 < H_0 < H_{C2}$	$\omega_{\pm} = \gamma(H_{C2} \pm H_0)$ (3.13)
	Noncollinear	$H_{C1} < H_0 < H_E^{\parallel}$	$\omega_1^{\parallel} = \gamma \left(\frac{2H_E H_E^{\perp}}{(H_E^{\parallel})^2} H_0^2 - 2H_E H_A \right)^{1/2}$ (3.14)
			$\{\omega_2^{\parallel} = 0\}$ (3.15)
	Parallel	$H_E^{\parallel} < H_0$	$\omega_1^{\parallel} = \gamma(H_0 + H_A)$ (3.16)
			$\{\omega_2^{\parallel} = \gamma(H_0 - H_E^{\parallel})\}$ (3.17)
$\mathbf{H}_0 \perp \mathbf{z}$	Noncollinear	$0 < H_0 < H_E^{\perp}$	$\omega_1^{\perp} = \gamma \left(H_{C2}^2 + \frac{H_E^{\parallel}}{H_E^{\perp}} H_0^2 \right)^{1/2}$ (3.18)
			$\omega_2^{\perp} = \gamma \left(H_{C2}^2 - \frac{H_A}{H_E^{\perp}} H_0^2 \right)^{1/2}$ (3.19)
	Parallel	$H_E^{\perp} < H_0$	$\omega_1^{\perp} = \gamma(H_0(H_0 - H_A))^{1/2}$ (3.20)
			$\{\omega_2^{\perp} = \gamma((H_0 - 2H_E)(H_0 - H_E))^{1/2}\}$ (3.21)

^a $\mathbf{z}$ is a unit vector parallel to the anisotropy axis.

^bField parameters: $H_E^{\parallel} = 2H_E - H_A$, $H_E^{\perp} = 2H_E + H_A$, $H_{C1} = \sqrt{H_A(2H_E - H_A)}$, $H_{C2} = \sqrt{H_A(2H_E + H_A)}$. For most antiferromagnets ($H_A \ll H_E$). Hence, $H_{C1} = H_{C2} = H_C = \sqrt{2H_E H_A}$ (Spin-flop field, $H_{spin-flop}$ as shown in figure 2.6), and $H_E^{\parallel} = H_E^{\perp} = 2H_E$ (Spin-flip field, $H_{spin-flip}$ as shown in figure 2.6)

^cBraces indicate the frequencies of modes that are not excited by a uniform ac magnetic field.

to the easy axis ($\mathbf{H}_A \parallel \mathbf{z}$). As seen from [table 3.1](#) and [figure 3.4](#), for $\mathbf{H}_0 \parallel \mathbf{z}$, an antiparallel (collinear) configuration is realized in weak fields. With increasing field strength, this state undergoes a transition to a non-collinear configuration via a spin-flop transition, which is a first-order phase transition where sublattices abruptly reorient at a critical field (H_c or $H_{\text{spin-flop}}$). Upon further increase of the magnetic field, the angle between $\mathbf{M}_1$ and $\mathbf{M}_2$ decreases continuously, and at $H_0 = H_E^\parallel$ a second-order spin-flip transition occurs, resulting in a parallel alignment of the two sublattice magnetizations along the field direction. In the case of $\mathbf{H}_0 \perp \mathbf{z}$, the angle between $\mathbf{M}_1$ and $\mathbf{M}_2$ decreases gradually with increasing H_0 favouring a non-collinear state. At $H_0 = H_E^\perp$, a transition from the non-collinear state to the parallel state takes place.

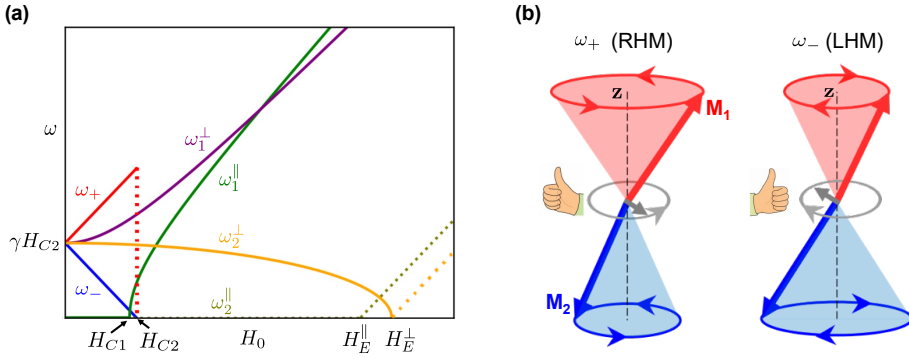


Figure 3.5: Antiferromagnetic eigenfrequencies and chiral modes in an easy-axis AF. (a) Eigenfrequencies of an AF with an easy-axis anisotropy. The subscripts $\parallel$ and $\perp$ correspond to $\mathbf{H}_0$ applied parallel and perpendicular to the easy axis, respectively. Dashed lines indicate modes with $M = 0$, which are not excited by a uniform rf field. (b) Circularly polarized chiral right-handed mode (RHM) and left-handed mode (LHM) in an easy-axis AF in the antiparallel ground state below the spin-flop transition, corresponding to the modes with frequencies ω_- and ω_+ , respectively.

In contrast to ferromagnetic eigenmodes, which always exhibit right-handed chirality due to the torque term $\mathbf{M} \times \mathbf{H}_{\text{eff}}$ fixing the sense of precession, antiferromagnets display richer and more intricate chiral dynamics. In the antiparallel configuration: i.e., below the spin-flop field ($0 < H_0 < H_{C2}$), the sublattice magnetizations $\mathbf{M}_1$ and $\mathbf{M}_2$, as well as the Néel vector $\mathbf{L}$, undergo circular precession that can possess either right-handed or left-handed chirality. These two chiral eigenmodes correspond to the ω_+ and ω_- branches, respectively (see [table 3.1](#)).

The right-handed mode (RHM) and left-handed mode (LHM) each carry a finite spin angular momentum along the easy axis, with opposite orientations^{96,183,188}. In the absence of an external magnetic field, these two modes are degenerate. A schematic illustration of the RHM and LHM, along with their field dependence on H_0 , is presented in [figure 3.5](#). For the uniaxial antiferromagnet MnF_2 , using the parameters $\mu_0 H_E = 53$ T and $\mu_0 H_A = 0.85$ T, the eigenfrequencies at zero external field ($H_0 = 0$ T) are $\omega_+ = \omega_- = \omega_0 = 266$ GHz¹⁸⁹. This corresponds to a relatively high resonance frequency characteristic of strong exchange coupling in AFs.

In the non-collinear state ($H_{C1} < H_0 < H_E^\parallel$), two distinct eigenmodes arise with eigenfrequencies $\omega_1^\parallel$ and $\omega_2^\parallel$, whose analytical expressions are given in [equation \(3.14\)](#) and [equation \(3.15\)](#), respectively (see [table 3.1](#)). The mode associated with $\omega_1^\parallel$ can be excited by an ac magnetic field transverse to the static field $\mathbf{H}_0$, and it is distinct from the two chiral modes $\omega_\pm$ discussed previously. The field dependence of this mode is illustrated in [figure 3.5a](#). The second mode, with eigenfrequency $\omega_2^\parallel$, cannot be excited by a spatially uniform ac magnetic field and is commonly referred to as the “soft mode”. This mode corresponds to an extremely slow collective rotation of $\mathbf{M}_1$ and $\mathbf{M}_2$ around the z-axis, with a frequency significantly smaller than those of the other eigenmodes in the system. Furthermore, in the parallel (spin-flip) state ($H_E^\parallel < H_0$), the corresponding eigenfrequencies $\omega_1^\parallel$ and $\omega_2^\parallel$ are given in [equation \(3.16\)](#) and [equation \(3.17\)](#), respectively (see [table 3.1](#)). In this regime, the first mode, described by [equation \(3.16\)](#), coincides with the uniform precessional mode of a uniaxial ferromagnet, whereas the second mode, [equation \(3.17\)](#), remains a soft mode.

For a magnetic field applied perpendicular to the easy axis, two distinct ground-state configurations are possible: the non-collinear and the parallel states (see [figure 3.4](#)). The corresponding eigenfrequencies for both configurations are summarized in [table 3.1](#), and their field dependence on H_0 is shown in [figure 3.5a](#). In the non-collinear state, the mode with eigenfrequency $\omega_1^\perp$ can be excited by an alternating magnetic field applied transverse to the static field $\mathbf{H}_0$. In this case, the net magnetization $\mathbf{M}_1 + \mathbf{M}_2$ precesses around $\mathbf{H}_0$ while maintaining a constant magnitude along its direction. The second mode in the non-collinear state, characterized by the eigenfrequency $\omega_2^\perp$, is instead excited by a longitudinal ac field (parallel to $\mathbf{H}_0$). Here, the net magnetization oscillates along the direction of $\mathbf{H}_0$ without developing any transverse ac component.

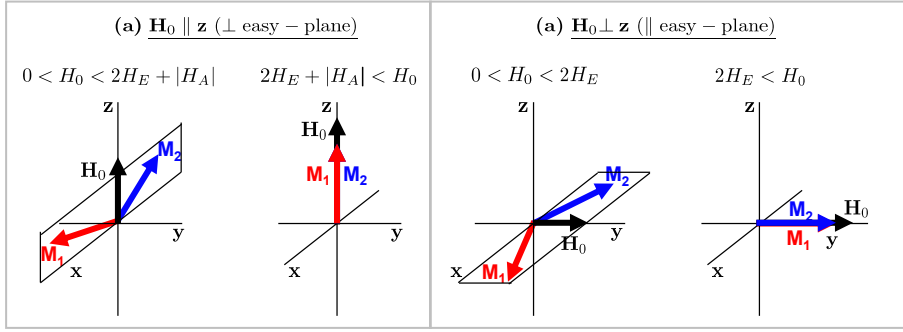


Figure 3.6: Ground-state configurations of a two-sublattice antiferromagnet with an easy-plane anisotropy for (a) a magnetic field applied perpendicular to the easy-plane ($\mathbf{H}_0 \parallel \mathbf{z}$) and (b) a magnetic field applied in the easy-plane direction ($\mathbf{H}_0 \perp \mathbf{z}$). $\mathbf{z}$ is unit vector perpendicular to the anisotropy plane. Adapted from [115].

3.3.2 Eigenmodes in easy-plane antiferromagnet

Let us now consider the case of an easy-plane AF, in which the easy directions of the sublattice magnetizations $\mathbf{M}_1$ and $\mathbf{M}_2$ lie within the basal or film plane. For this system, two distinct ground-state configuration, non-collinear and parallel can exist for both orientations of the external magnetic field $\mathbf{H}_0$: when it is applied perpendicular to the basal plane (i.e., $\mathbf{H}_0 \parallel \mathbf{z}$) and when it is applied within the plane (i.e., $\mathbf{H}_0 \perp \mathbf{z}$). The corresponding ground-state configurations are illustrated in figure 3.6, while the analytical solutions for the associated eigenfrequencies are summarized in table 3.2¹¹⁵.

As can be seen by comparing figure 3.4 and figure 3.6, the non-collinear and parallel ground states for $\mathbf{H}_0 \parallel \mathbf{z}$, as well as the parallel ground state for $\mathbf{H}_0 \perp \mathbf{z}$, in an easy-plane AF are structurally identical to those of an easy-axis antiferromagnet under the same field orientations. Consequently, the corresponding expressions for the eigenfrequencies are formally the same in both cases. The key distinction lies in the sign of the anisotropy field H_A , which is negative for the easy-plane AF. This change in sign leads to characteristic differences in the field dependence of the eigenmodes with respect to H_0 . The field dependencies of all eigenmodes in the easy-plane case are shown in figure 3.7a.

In easy-plane AFs, a distinct eigenfrequency arises in the non-collinear state for the case $\mathbf{H}_0 \perp \mathbf{z}$. This corresponds to the mode $\omega_1^\perp$ (see equation (3.24) in

Table 3.2: Eigenfrequencies of a two-sublattice AF with easy-plane anisotropy ($H_A < 0$) at $T = 0K$. Reproduced from [115].

Field direction ^a	Ground state	Field range	Eigenfrequencies
$\mathbf{H}_0 \parallel \mathbf{z}$	Noncollinear	$0 < H_0 < 2H_E + H_A $	$\omega_1^{\parallel} = \gamma \left(\frac{2H_E H_E^{\perp}}{(H_E^{\parallel})^2} H_0^2 - 2H_E H_A \right)^{1/2} \quad (3.22)$
	Parallel	$2H_E + H_A < H_0$	$\omega_1^{\parallel} = \gamma(H_0 + H_A) \quad (3.23)$
$\mathbf{H}_0 \perp \mathbf{z}$	Noncollinear	$0 < H_0 < 2H_E$	$\omega_1^{\perp} = \gamma H_0 \left(1 + \frac{ H_A }{2H_E} \right)^{1/2} \quad (3.24)$
			$\omega_2^{\perp} = \gamma \left(2H_E H_A - \frac{ H_A }{2H_E} H_0^2 \right)^{1/2} \quad (3.25)$
	Parallel	$2H_E < H_0$	$\omega_1^{\perp} = \gamma (H_0(H_0 - H_A))^{1/2} \quad (3.26)$

^a $\mathbf{z}$ is a unit vector perpendicular to the anisotropy plane.

table 3.2), whose field dependence is shown in figure 3.7a. The corresponding character of the oscillations of $\mathbf{M}_1$ and $\mathbf{M}_2$ in this state is illustrated schematically in figure 3.7b. This mode can be excited by the transverse component of an ac magnetic field (with respect to $\mathbf{H}_0$), directed along the y -axis. Its frequency approaches zero as $H_0 \rightarrow 0$, reflecting its acoustic nature. Physically, this mode, often referred to as the acoustic mode (AM), can be understood as a precession of the sublattice magnetizations $\mathbf{M}_1$ and $\mathbf{M}_2$ in which the angle between them remains constant. Consequently, the net magnetization ($\mathbf{m}$) oscillates around the direction of the applied magnetic field $\mathbf{H}_0$ without any change in its magnitude.

The second mode, $\omega_2^{\perp}$, in the non-collinear state is excited by a longitudinal ac magnetic field (parallel to $\mathbf{H}_0$) (see equation (3.25) in table 3.2). The field dependence of this mode, along with the corresponding oscillation characteristics, is illustrated in figure 3.7. This mode is commonly referred to as the optical mode (OM). It can be understood as a collective oscillation of the sublattice

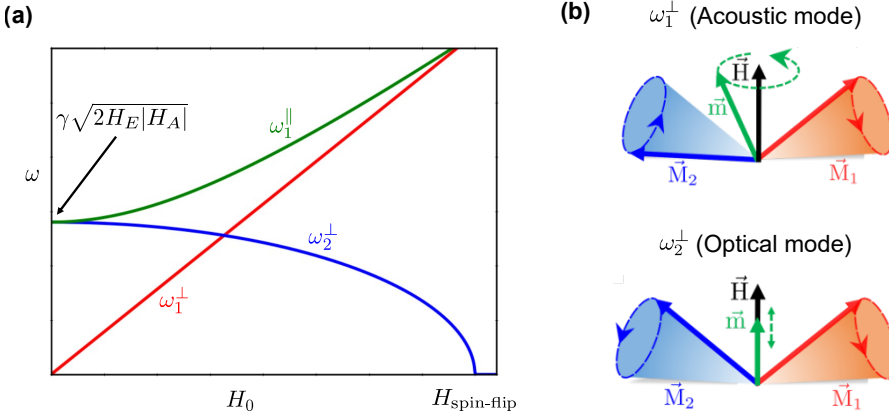


Figure 3.7: (a) Eigenfrequencies of an antiferromagnet with an easy-plane anisotropy. The subscripts $\parallel$ and $\perp$ correspond to $\mathbf{H}_0$ applied parallel and perpendicular to the $\mathbf{z}$ axis, respectively where $\mathbf{z}$ is a unit vector perpendicular to the anisotropy plane. (b) Precession of the sublattice-magnetization ($\mathbf{M}_1$ and $\mathbf{M}_2$) and total-magnetization ($\mathbf{m}$) vectors of acoustic and optical eigenmodes of an easy-plane antiferromagnet corresponding to $\omega_1^{\perp}$ and $\omega_2^{\perp}$ eigenmode branches from (a).

magnetizations $\mathbf{M}_1$ and $\mathbf{M}_2$, in which their relative orientations vary such that the amplitude of the net magnetization ($\mathbf{m}$) oscillates along the direction of the applied field $\mathbf{H}_0$, without developing a transverse component.

The OM typically exhibits a much higher resonance frequency compared to the AM and remains finite even at zero magnetic field. This behavior reflects the fact that the mode is predominantly governed by the strong internal exchange and anisotropy fields rather than by the external field. Both the acoustic and optical eigenmodes are linearly polarized as a consequence of the highly elliptical precession of the sublattice magnetizations $\mathbf{M}_1$ and $\mathbf{M}_2$, which possess opposite chiralities^{96,188,190,191}. Owing to this symmetry, these modes carry vanishing net spin angular momentum on average^{96,188,192}. Another distinctive feature of the easy-plane antiferromagnetic resonance spectrum is the crossover of the acoustic and optical eigenmode branches, corresponding to a degeneracy of the modes at a critical field $H_0 = H_{\text{cross-over}}$. This degeneracy is lifted when the applied field $\mathbf{H}_0$ is tilted out of the easy plane, leading to anti-crossing and a splitting of the previously coincident modes.

3.4 Spin waves or magnons

The solutions of the LLG equation that describe the collective precessional motion of magnetic moments are known as *spin waves* (SWs)^{98,115,179,193,194}. The quanta of these waves, referred to as *magnons*, are bosonic quasiparticles corresponding to a spatially delocalized flip of a single electron spin. Each magnon carries a spin angular momentum of $S_m = \hbar$, an energy $\epsilon = \hbar\omega$, and a wavevector magnitude $|q| = 2\pi/\lambda$, where λ is the spin-wave wavelength^{98,195–197}. The existence of SWs was first predicted by F. Bloch in 1930, when he calculated the SW spectrum for a chain of ferromagnetically aligned spins¹⁹⁵. The first direct experimental observation of SWs came in 1946 with the discovery of FMR by Griffiths¹⁷⁵, followed by the discovery of AFMR in 1952 by Ubbink and colleagues¹⁸⁵. In the previous sections discussing FMR and AFMR, only a *homogeneous precession* of the magnetization $\mathbf{M}$ or the Néel order parameter $\mathbf{L}$ throughout the entire sample volume was considered. In such a uniform mode, all magnetic moments precess at the same frequency and with the same phase. This corresponds to a SW with infinite wavelength, or equivalently, a zero wavevector ($q = 0$). Since the quanta of SWs are magnons, the FMR or AFMR mode represents a zero-wavevector magnon¹⁷⁸.

Higher-order excitations give rise to modes in which the magnetization precesses at the same frequency but with a spatially varying phase. In this case, neighboring magnetic moments precess slightly out of phase, forming a *propagating SW* with a finite wavelength ($\lambda < \infty$) and nonzero wavevector ($q \neq 0$)^{115,179,193,194}. Such a wave can be modeled as a chain of exchange-coupled magnetic moments precessing around their equilibrium directions with a continuously varying phase, as schematically illustrated in Figure 3.8a. For a one-dimensional chain of spins coupled ferromagnetically via nearest-neighbor exchange J , the SW dispersion relation is given by

$$\hbar\omega_q = \hbar\gamma\mu_0 H_0 + 4JS(1 - \cos(qa)), \quad (3.27)$$

where $\hbar$ is the reduced Planck constant, ω_q is the angular frequency of the SW with wavevector q , γ is the gyromagnetic ratio, μ_0 is the vacuum permeability, H_0 is the external magnetic field applied along the quantization axis, J is the exchange constant describing the strength of the nearest-neighbor ferromagnetic coupling, S is the spin quantum number of each magnetic moment, and a is the lattice

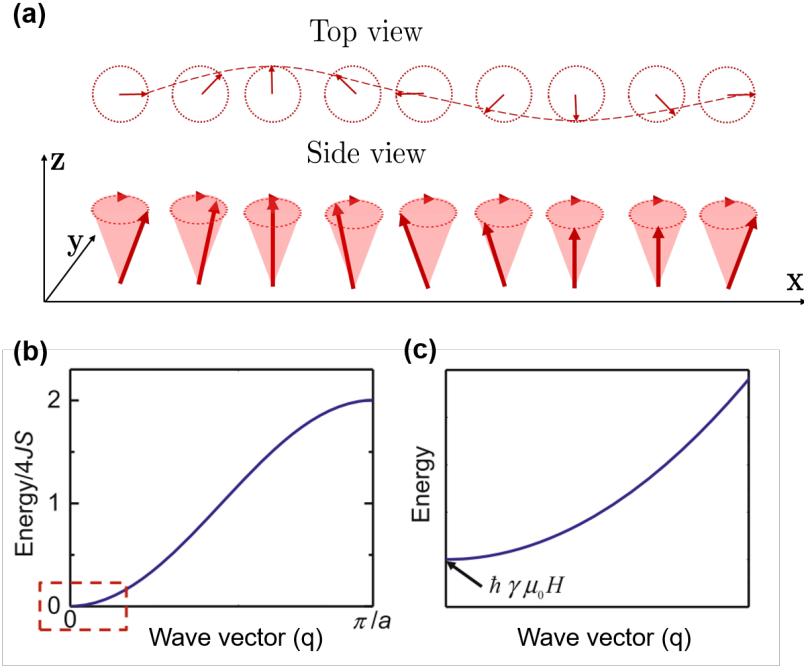


Figure 3.8: (a) Illustration of a SW propagating in the $+x$ -direction along a one-dimensional chain of ferromagnetically coupled classical spins. The phase of precession varies continuously along the chain, and the distance between the two spins at the ends with no phase difference corresponds to one wavelength λ . (b) Dispersion relation for magnons in a one-dimensional linear chain over the full Brillouin zone. (c) Magnified view near the origin ($qa \ll 1$), where the dispersion exhibits a quadratic dependence on q and the Zeeman energy gap $\hbar \gamma \mu_0 H_0$ is visible at $q = 0$.

spacing between neighboring spins. In the long-wavelength limit ($qa \ll 1$), the cosine term can be expanded to second order, yielding a quadratic dispersion

$$\hbar \omega_q \approx \hbar \gamma \mu_0 H_0 + 2JSa^2 q^2. \quad (3.28)$$

This clearly shows that the SW energy increases quadratically with wavevector for small q and that the minimum energy at $q = 0$ is given by $\hbar \omega_0 = \hbar \gamma \mu_0 H_0$, corresponding to all spins precessing in phase, where the exchange interaction does not contribute to the energy^{179,194}. As the wavevector increases, the phase difference between neighboring spins grows, increasing the exchange energy

contribution. The full dispersion relation (3.27) and its quadratic long-wavelength approximation (3.28) are illustrated in Figures 3.8b and 3.8c, highlighting both the quadratic dependence at small q and the energy gap $\hbar\gamma\mu_0 H_a$ (at non-zero magnetic anisotropy) at $q = 0$.

Likewise to the case of chain of ferromagnetically aligned spins, the spectrum of a magnon in an isotropic FM is generally gapless at $H_0 = 0$ (neglecting anisotropy). However, the magnon spectrum exhibits an energy gap, meaning its energy ε_0 is finite at $\mathbf{q} = 0$ which is determined by intrinsic material parameters such as the effective exchange fields H_{ex} , magnetic anisotropy field H_A , and dipole-dipole interactions, as well as by an externally applied magnetic field H_0 . In the case of an isotropic Heisenberg FM, the SW excitations can be described by a quadratic dispersion relation in the long-wavelength limit (low q -values). The spin-wave frequency is given by

$$\omega_k = \gamma\mu_0 (H_0 + Dq^2), \quad (3.29)$$

where the term D denotes the *exchange stiffness parameter*, defined as $D = \frac{2JSa^2}{g\mu_B\mu_0}$ with g the Landé g -factor, and μ_B the Bohr magneton^{193,194}. At finite damping, the dynamics of the spin precession are affected by the Gilbert damping parameter α . For $\alpha \ll 1$, the damping can be treated as a small perturbation to the ideal precession. When dipolar interactions are included, as in thin films or confined geometries, the dispersion becomes anisotropic and depends on the propagation direction relative to the magnetization. This leads to various magnetostatic SW modes, such as the Damon–Eshbach and backward-volume modes, each with distinct dispersion characteristics^{193,194,198}.

A simple FM exhibits a single branch in its SW spectrum, corresponding to the coherent precession of all spins in the lattice. In contrast, multi-sublattice systems such as AFs and ferrimagnetic materials possess multiple magnon branches arising from the coupled dynamics of the individual magnetic sublattices^{187,194}. The number and nature of these branches depend on the mutual orientation of the sublattice magnetizations. In an isotropic AF without any anisotropy magnon dispersion at smaller $\mathbf{q}$ is usually linear and is quite distinct from the quadratic dispersion of the FMs. In two-sublattice AF systems, distinct dispersion relations are obtained for different magnetic configurations, including antiparallel, non-collinear, and parallel alignments of the sublattice magnetizations, both in

easy-axis and easy-plane AFs, as discussed in the previous section. Furthermore, in easy-plane antiferromagnets, the magnon spectrum typically consists of two fundamental modes: the acoustic and optical mode SW branches. The AM corresponds to an in-phase precession of the sublattice magnetizations and exhibits a gapless (or nearly gapless) linear dispersion at low wave vectors, whereas the OM originates from an out-of-phase precession and is characterized by a finite frequency gap at $\mathbf{q} = 0$. These multiple dispersions reflect the rich collective dynamics of multi-sublattice magnetic systems and provide a basis for understanding the complex SW behavior observed experimentally in antiferromagnetic and ferrimagnetic materials. Figure 3.9 illustrates the spin-wave dispersion relation for an easy-axis AF in the antiparallel configuration at zero applied magnetic field. In this configuration, the two degenerate circularly polarized SW mode frequencies can be expressed as

$$\omega_{\alpha\mathbf{k}} = \omega_{\mathbf{k}} + \gamma H_0, \quad \omega_{\beta\mathbf{k}} = \omega_{\mathbf{k}} - \gamma H_0, \quad (3.30)$$

$$\omega_{\mathbf{k}} = \gamma \sqrt{2H_E H_A + H_A^2 + H_E^2 (1 - \gamma_{\mathbf{k}}^2)}, \quad (3.31)$$

where H_E denotes the inter-sublattice exchange field and H_A the uniaxial anisotropy field. The geometric (or structural) factor $\gamma_{\mathbf{k}}$ depends on the crystal lattice symmetry and can be written as

$$\gamma_{\mathbf{k}} = \cos\left(\frac{k_x a}{2}\right) \cos\left(\frac{k_y a}{2}\right) \cos\left(\frac{k_z c}{2}\right), \quad (3.32)$$

where a and c are the lattice parameters along the crystallographic a - and c -axes, respectively^{187,193,194}. When a finite magnetic field H_0 is applied, the degeneracy of the magnon modes is lifted, leading to a splitting of the two SW branches, $\omega_{\alpha\mathbf{k}}$ and $\omega_{\beta\mathbf{k}}$.

The field of *magnonics*, wherein magnons are harnessed to transmit and process information at the nanoscale, stands as a promising alternative to conventional microelectronics based on moving charges. SWs can serve as data carriers, offering several intrinsic advantages over electrons in circuits. For instance, SW frequencies typically start in the microwave (GHz) regime and may extend into the terahertz (THz) range, thus supporting ultrafast operation and high clock rates^{96,188,198,199}. Even at microwave frequencies, their associated wavelengths can

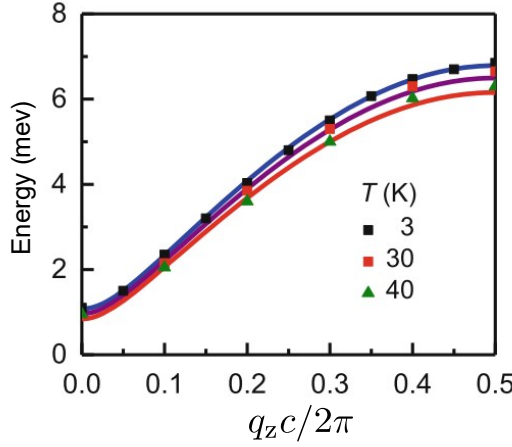


Figure 3.9: Magnon dispersion relations for easy-axis antiferromagnet MnF_2 in the absence of an externally applied magnetic field, with the wave vector oriented along the crystallographic c -axis. The symbols denote the experimental data obtained from the inelastic neutron scattering measurements at three different temperatures. The solid lines represent the corresponding magnon dispersion curves calculated for the same temperatures. Reproduced from [194].

shrink to the nanometre scale, allowing for devices that are substantially smaller than their electronic or photonic counterparts²⁰⁰. Moreover, the characteristics of SWs, such as dispersion, amplitude, and phase, can be widely tuned through the choice of magnetic material, sample geometry, and external magnetic field²⁰¹. Nonlinear SW phenomena further open additional degrees of freedom for data processing, enabling functionalities such as phase encoding, interference, and wave multiplexing²⁰².

Although these attributes provide strong motivation for “wave-based computing,” existing ferromagnet-based magnonic devices face several intrinsic challenges. Chief among these are the inefficient excitation of SWs at the nanoscale and the relatively short propagation lengths caused by magnetic damping in thin ferromagnetic films¹⁹⁹. SOT techniques have been shown to partially overcome these issues by compensating damping and even generating auto-oscillations in ferromagnetic nanostructures^{79–84}, which is the discussion of the next section. Nevertheless, ferromagnetic systems remain constrained by their GHz-range dynamics, stray magnetic fields, and limited control over mode chirality and multi-

iband operation. These limitations impose severe restrictions on the scalability and coherence of FM-based magnonic circuits.

In this context, AF materials have recently emerged as a compelling alternative. The dynamics of AFs are governed by strong exchange coupling and magnetic anisotropy, leading to qualitatively distinct behaviour compared to FMs. As emphasised in the recent review by Han et al.¹⁸⁸, AFs provide a pathway toward coherent SW dynamics, phase-correlated spin transport, and ultrafast SOT effects. Their characteristic eigenfrequencies often lie in the THz regime, offering a natural platform for ultrafast operation^{3,5}. The absence of net magnetisation eliminates stray magnetic fields, thereby reducing magnetic cross-talk and enhancing integration density^{3,5}. Furthermore, experiments suggest that spin currents in AFs can propagate as coherent evanescent SWs, supporting phase-stable information transfer over nanoscale distances⁹⁰. Owing to their multiband magnonic spectra, AFs also support rich mode structures and offer new opportunities for control via SOTs and ultrafast optical excitations^{5,188}.

3.5 SOT-driven self-oscillations

So far, we have discussed the excitation of magnetization dynamics, namely FMR and AFMR, driven by an external oscillating rf magnetic field. In [chapter II](#), we introduced current-induced SOTs as a fully electrical mechanism to excite magnetic order. As will be demonstrated as results of this thesis in [chapter V](#) and [chapter VI](#), SOTs can effectively drive coherent oscillations of the Néel order in SAFs as well as excite intrinsic antiferromagnetic eigenmodes.

In this section, we focus on one of the most important and promising applications of SOTs: *SOT-driven self-oscillations* in both FMs and AFs. The concept of current-induced self-oscillation was first proposed by Slonczewski in 1989 in the context of STT¹⁰⁴. He predicted that, depending on the current polarity, STT can either enhance the effective Gilbert damping or compensate it entirely, thereby enabling spontaneous steady-state oscillations. Since this original prediction, numerous experimental demonstrations of STT-driven self-oscillations in spin valves and MTJs^{73–78}, as well as SOT-driven self-oscillations^{79–84} in HM/FM heterostructures, have been reported.

STT-based oscillators are commonly referred to as *spin-torque nano-oscillators* (STNOs), whereas SOT-based oscillators are known as *spin Hall nano-oscillators* (SHNOs), emphasizing the role of the SHE as the underlying source of SOT. The qualifier “nano” reflects the device geometries typically employed, such as nano-contacts, nano-pillars, or nano-constrictions, which enable the high current densities required to overcome intrinsic damping and trigger magnetic self-oscillations⁷³.

To describe the effect of SOTs on the magnetization dynamics, we consider the LLG equation augmented by FL and DL torque terms:

$$\frac{d\mathbf{M}}{dt} = -\gamma \mathbf{M} \times \mu_0 \mathbf{H}_{\text{eff}} + \frac{\alpha}{M_{\text{sat}}} \left(\mathbf{M} \times \frac{d\mathbf{M}}{dt} \right) + \tau_{\text{FL}} \mathbf{M} \times \mathbf{P} + \tau_{\text{DL}} \mathbf{M} \times (\mathbf{M} \times \mathbf{P}), \quad (3.33)$$

where the last two terms represent the FL and DL SOT components, respectively. Here, $\mathbf{P}$ denotes the unit vector along the spin polarization direction of the injected spin current. For a Pt layer acting as the HM layer on top of a FM, and assuming a positive SHA of Pt, a charge current flowing along $+\hat{\mathbf{x}}$ generates a spin current flowing along $-\hat{\mathbf{z}}$ into the FM with spin polarization oriented along $-\hat{\mathbf{y}}$. It is convenient to introduce two effective magnetic fields, $\mathbf{H}_{\text{DL}}$ and $\mathbf{H}_{\text{FL}}$, corresponding to the DL and FL SOTs respectively. The torques can be expressed in terms of these effective fields via $\mathbf{T}_{\text{DL,FL}} = \mathbf{M} \times \mathbf{H}_{\text{DL,FL}}$. With $\mathbf{P}$ denoting the unit vector along the spin polarization direction, the effective fields take the form $\mathbf{H}_{\text{DL}} = H_{\text{DL}} \mathbf{M} \times \mathbf{P}$ and $\mathbf{H}_{\text{FL}} = H_{\text{FL}} \mathbf{P}$. If the torques are written for the normalized magnetization vector, the field amplitudes are directly given by $H_{\text{DL}} = \tau_{\text{DL}}$ and $H_{\text{FL}} = \tau_{\text{FL}}$. To quantify the torque strength, it is useful to define the SOT efficiencies

$$\xi_j^{\text{DL,FL}} = \frac{2e}{\hbar} \frac{M_s t_F}{j_c} H_{\text{DL,FL}}, \quad (3.34)$$

where M_s is the saturation magnetization, t_F the ferromagnetic layer thickness, and j_c the applied charge current density. The parameters $\xi_j^{\text{DL,FL}}$ represent the ratio of the effective spin current absorbed by the FM to the injected charge current in the adjacent nonmagnetic layer and can thus be interpreted as effective SHAs for a given HM/FM bilayer system. The qualitative actions of FL and DL SOTs on the dynamic magnetization are illustrated in [figure 3.10](#). The DL torque τ_{DL} can either enhance or compensate the intrinsic Gilbert damping depending on the current polarity; thus, it can either suppress precession or enable sustained self-

oscillations. The field-like torque τ_{FL} , having the same symmetry as an effective magnetic field, modifies the precession frequency by shifting the total effective field.

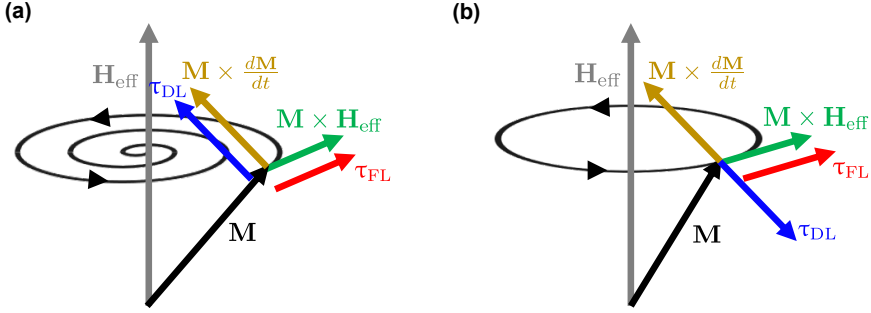


Figure 3.10: (a) Damped precessional motion of the magnetization $\mathbf{M}$ under the action of SOTs. Here, the Gilbert damping torque and the damping-like SOT act in the same direction, thereby enhancing the relaxation of $\mathbf{M}$ toward the effective field $\mathbf{H}_{eff}$. (b) Self-oscillation of the magnetization $\mathbf{M}$ arising from the compensation of the Gilbert damping torque by an oppositely oriented damping-like SOT, enabling sustained precession.

To induce persistent oscillations or magnetization switching, the spin current density delivered to the ferromagnetic layer must exceed a threshold sufficient to counteract intrinsic damping. Achieving such large spin currents motivates the use of specialized nano-geometries in SHNOs⁷³. Notably, Demidov *et al.* proposed nanoconstriction-based SHNOs, in which a narrow constriction confines the charge-current path, resulting in a substantial enhancement of the local current density and enabling self-oscillations even at moderate applied currents⁸². Compared to earlier designs, nanoconstriction SHNOs offer advantages such as simpler fabrication, reduced operating currents, and improved localization of the oscillation mode and in this thesis we will be exploring this nano-constriction geometry in or SAF systems.

The advantages of SOT-driven nano-oscillators have been outlined in [chapter I](#), together with their known limitations and the alternative approaches explored in this thesis. The current state-of-the-art devices rely predominantly on conventional FMs, which suffer from several drawbacks: relatively slow magnetic dynamics, a limited number of accessible oscillation modes, a restricted range of current-tunable frequencies, the inability to control the chirality of the self-oscillations via current, and, in general, the inherent disadvantages associated

with FMs for memory integration. In contrast, AFs-based SOT-driven nano-oscillators have the potential to overcome all of these limitations. However, only a handful of theoretical predictions currently exist, indicating a promising yet largely unexplored research direction.

These considerations define the motivation for investigating the SOT-driven dynamics of SAFs in this thesis. The next chapter begins this exploration with the first experimental observation of a SOT-driven, non-volatile antiferromagnetic memristor based on a SAF system.



IV

SOT control and readout of Néel order reversal in SAFs

Building on the concepts of SOTs and the unique advantages of AFs introduced in previous chapters, this chapter demonstrates the controlled electrical manipulation and detection of the Néel-order in a SAF with PMA. SOT drives DW motion to “write” reconfigurable, nonvolatile states, while nonlinear magnetotransport provides a precise “readout” of the antiferromagnetic domains. The absence of long-range stray fields ensures stable, reproducible, and pulse-driven control of domain configurations. These results highlight the potential of SAF as a versatile nonvolatile memory mimicking an artificial synapse, with applications in neuromorphic computing, where the reconfigurable synaptic weight can be encoded in the ratio of reversed antiferromagnetic domains.

Keywords: Non-linear magnetotransport, spin-orbit torques (SOTs), memristor, neuromorphic spintronics, antiferromagnetic spintronics.

This chapter is based on: Godinho, J., **Rout, P.K.**, et al. Antiferromagnetic domain wall memory with neuromorphic functionality. *npj Spintronics* 2, 39 (2024)²⁰³.

4.1 Introduction

Up to this point it has become evident that AFs hold exceptional promise for next-generation spintronic technologies owing to their unique properties such as the absence of stray magnetic fields, robustness against external perturbations, ultrafast spin dynamics, and the potential for high-density device integration^{2–6,96}. These characteristics make AFs ideal candidates for realizing energy-efficient and scalable spintronic devices. However, two major challenges continue to limit their practical implementation: the efficient electrical manipulation and reliable electrical detection of the antiferromagnetic order^{3,5}. Unlike ferromagnets, the compensated spin structure in AFs leads to a vanishing net magnetization, which makes it inherently difficult to control and probe their magnetic state using conventional techniques. Overcoming these limitations is essential for developing CMOS-compatible memory and logic devices that can fully exploit the speed, stability, and scalability offered by antiferromagnetic materials.

Recently, neuromorphic computing has emerged as a promising approach to achieve energy-efficient processing of large-scale data by emulating the functionality of the human brain (see [section 1.3](#))^{50,51}. In artificial neural networks, threshold-activated neurons are interconnected through synapses whose analog weights are adjusted during learning. Such architectures underpin a wide range of modern applications, including pattern or facial recognition, language translation, and GPT-like large language models. However, as computational demands grow, their energy consumption becomes unsustainable. To overcome this limitation, hardware-based neuromorphic concepts are being developed, exploiting resistive phase transitions or nonlinear magnetization dynamics to replicate neuronal and synaptic functions^{52–55,57,60,61}. Most existing memristive implementations, however, rely on ionic motion or thermally induced structural changes, which often limit their reliability and scalability^{64–67}.

One promising route toward energy-efficient artificial synapses is using magnetic systems, where the synaptic weight is encoded in the number or position of magnetic textures such as skyrmions^{71,72} or DWs^{69,70,204}. A well-known example of a non-volatile magnetic synapse is the MTJ, where current-induced DW motion driven by STT enables partial magnetization reversal ([figure 4.1a](#))⁷⁰. However, ferromagnetic systems inherently produce long-range stray fields that can perturb

nearby nanostructures during switching. These interactions, along with thermal fluctuations, lead to stochastic DW pinning and depinning (figure 4.1b), hindering precise and reproducible control of analog synaptic weights.

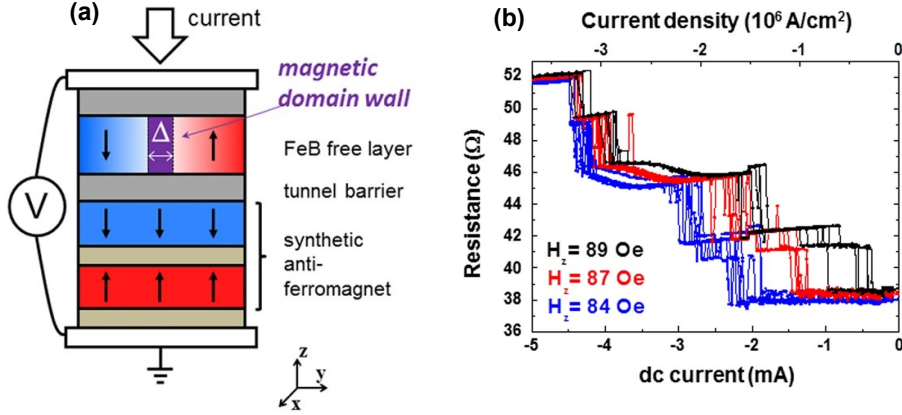


Figure 4.1: Spin-torque memristor. (a) Schematic of an MgO-based MTJ featuring a DW in the FeB free layer, where Δ represents the DW width. (b) Resistive variations measured during current-driven DW depinning from a specific pinning center (resistance value of 38 Ω) under different external magnetic fields in the range of 78–92 Oe. Reproduced from [70].

AFs present a promising alternative for neuromorphic spintronics, addressing several limitations inherent to FM-based devices. Recent advancements have demonstrated the development of antiferromagnetic synaptic elements, wherein electrical or optical excitation heats the antiferromagnetic film near its Néel temperature^{205,206}. This process fragments large antiferromagnetic domains into multiple submicron domains, inducing a high-resistance multidomain state. Upon cooling, the system reverts to a low-resistance, few-domain configuration, enabling spike-timing-dependent plasticity in neural networks, though without non-volatile functionality. This instability stands in marked contrast to the robust, non-volatile SOT-induced 180° Néel vector reversal previously demonstrated in CuMnAs^{15,207}, a mechanism we now seek to utilize within a synthetic antiferromagnetic system possessing nominally compensated magnetic moments. Storing information in domains with reversed Néel order is more stable than the quench switching generating a given ratio of orthogonally oriented domains.

In this work, we employ a compensated SAF system with PMA, where non-volatile analog information is encoded in stable antiferromagnetic domains with Néel vectors oriented along the magnetic easy axis. The combined $\mathcal{PT}$ -symmetry of the SAF multilayer enables SOT-driven 180° DW motion to tune the proportion of reversed domains, allowing access to analog information via nonlinear transport measurements. We observe highly reproducible writing, which we attribute to the absence of magnetic stray fields in the compensated antiferromagnetic system.

4.2 Experimental Methods

4.2.1 Sample and device fabrication

For the realization of non-volatile antiferromagnetic memory, we employ a nominally compensated SAF with the following stack: Si/SiO₂(sub.)/Ta(1.5nm)/Pt(5nm)/FM2//Ir(0.6nm)/FMI/Pt(4nm)/Ru(3nm). Each ferromagnetic layer, FMI and FM2, consists of the following Co/Ni multilayers, [Co(0.25nm)/Ni(0.65nm)]_{x3}/Co(0.25nm), exhibiting a PMA. As shown in [figure 4.2a](#), the two ferromagnetic layers are antiferromagnetically coupled via the RKKY interaction through the thin Ir layer. The stack is sandwiched between nonmagnetic Pt layers, where current-induced spin accumulations of opposite polarity are generated at the interfaces via the SHE and Rashba–Edelstein effects, mimicking the staggered Néel SOT observed in $\mathcal{PT}$ -symmetric AFs such as CuMnAs and Mn₂Au^{15,37}.

The SAF stack was fabricated by dc magnetron sputtering in an ultrahigh vacuum ATC 2200 system (AJA International Inc.) under 3×10^{-3} mbar Ar. Si substrates with 100 nm thermally oxidized SiO₂ were used. A 1.5 nm Ta adhesion layer was deposited first, followed by a 5 nm Pt seed layer to enhance perpendicular magnetic anisotropy. The multilayer was grown at room temperature with alternating source shutters, and the sample holder rotated at 60 rpm to ensure uniform deposition. Finally, a 3 nm Ru capping layer was added to prevent oxidation, with layer thicknesses controlled via deposition time.

For transport measurements, Hall-bar devices were fabricated via electron-beam lithography using positive PMMA resist. Device geometries were defined by Ar ion milling, followed by deposition of Cr/Au (5/80 nm) electrical contacts

using a lift-off process. Figure 4.2b shows a scanning electron micrograph of a Hall-bar device alongside a schematic of the measurement setup.

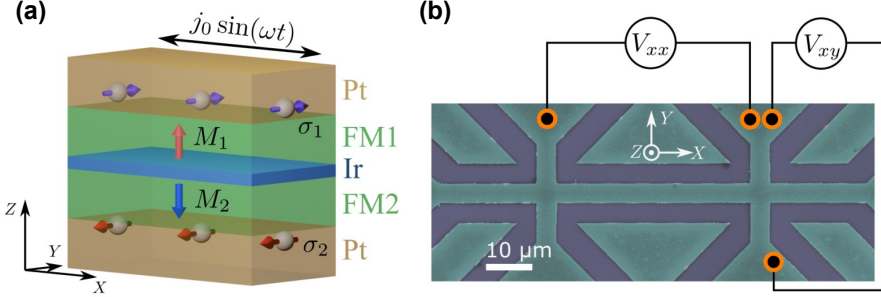


Figure 4.2: (a) Stack structure of the perpendicularly magnetized SAF. Due to the SHE, spin accumulations (σ_i) with opposite orientations are generated at the Pt/FM interfaces, mimicking the staggered Néel SOT. (b) Scanning electron micrograph of the Hall-bar device, alongside a schematic of the measurement setup.

4.2.2 Detection of Néel order by non-linear magnetoresistance

For the detection of antiferromagnetic Néel order, an odd effect such as the AHE is required to distinguish oppositely oriented Néel vectors. As discussed in [section 2.4.3](#), the AHE response vanishes in $\mathcal{PT}$ -symmetric AFs. Similarly, in a fully compensated SAF, the AHE signal is zero. However, linear AHE has been reported in intentionally uncompensated SAF systems, where a net magnetic moment breaks $\mathcal{PT}$ -symmetry and enables differentiation between AF states^{208,209}. Moreover, the linear AMR response is even in magnetization and therefore cannot detect a 180° reversal of the Néel vector.

The proposed approach relies on the measurement of second-order magnetotransport effects. The underlying microscopic mechanism arises from the interplay between magnetoresistance and current-induced fields. In the presence of such fields, the resistance becomes current-dependent as magnetic moments are displaced from their equilibrium positions. The first-order, or linear, response is described by $E_i = \rho_{ij}j_j$, where ρ_{ij} represents the current-independent resistivity. Higher-order terms capture the current dependence of the resistivity; here, we consider contributions up to the second order. For a second-order response to occur, inversion symmetry must be broken in the system. The second-order

transport equation can be expressed as $E_i = \xi_{ijk} j_j j_k$, where both the electric field E_i and current density j_i are odd under inversion symmetry $\mathcal{P}$. Consequently, the $\mathcal{P}$ symmetry must be broken for ξ_{ijk} to be nonzero. Figure 4.3 illustrates a simple mechanism responsible for the emergence of the second-order term. The direction in which the AF moments tilt away from equilibrium depends on the orientation of the Néel vector. In $\mathcal{PT}$ -symmetric AF CuMnAs, an applied electric current induces a staggered spin polarization that generates spin-orbit fields $\mathbf{H}_{\text{SO}}^1$ and $\mathbf{H}_{\text{SO}}^2$, satisfying $\mathbf{H}_{\text{SO}}^1 = -\mathbf{H}_{\text{SO}}^2$. These fields act on the corresponding sublattice magnetizations $\mathbf{M}_1$ and $\mathbf{M}_2$. The spin-orbit fields are oriented perpendicular to the applied current direction, and their magnitude scales linearly with the current density j . In equilibrium, $\mathbf{M}_1 = -\mathbf{M}_2$. The associated SOTs, $\mathbf{H}_{\text{SO}}^1 \times \mathbf{M}_1$ and $\mathbf{H}_{\text{SO}}^2 \times \mathbf{M}_2$, cant the sublattice magnetizations away from their perfectly antiparallel configuration. The resulting exchange torques subsequently drive a coherent rotation of $\mathbf{M}_1$ and $\mathbf{M}_2$ within the basal plane of CuMnAs toward the directions of the spin-orbit fields $\mathbf{H}_{\text{SO}}^{1,2}$. The detection scheme exploits the fact that both the SOTs and the resulting exchange torques reverse sign upon reversal of the sublattice magnetizations. Consequently, the reversed Néel vector is deflected in the opposite sense (clockwise versus counterclockwise). In combination with magnetization-even AMR, this leads to a second-order magnetoresistance response that is odd and is generally different for the two reversed states, thereby enabling purely electrical detection of the Néel vector reversal¹⁵. However, as will be shown in the results section, it is not the FL spin-orbit field but rather the DL spin-orbit field that is responsible for the second-order magnetoresistance signal observed in SAF with PMA when an external magnetic field is applied perpendicular to the easy-axis anisotropy direction.

The displacements of the AF moments from their equilibrium orientations induced by SOTs are typically very small, resulting in a correspondingly weak second-order magnetoresistance signal. To enhance the detection sensitivity, we employ a homodyne measurement technique. The experiments are carried out on a $50 \times 5 \mu\text{m}$ Hall bar device (figure 4.2b), with a low-frequency (163 Hz) probing current applied through the sample. Using lock-in amplifiers, we simultaneously record four signals: the longitudinal and transverse voltages, along with their first- and second-harmonic responses. From Ohm's law, the voltage can be expressed as:

$$V = R(\theta)I. \quad (4.1)$$

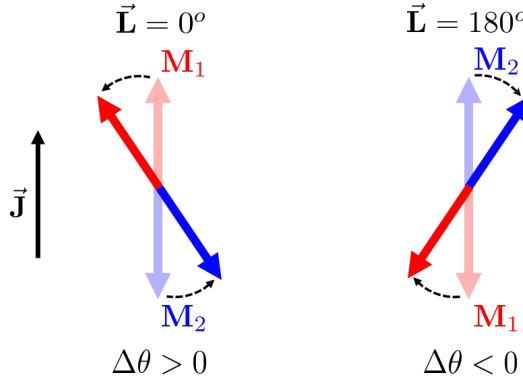


Figure 4.3: Illustration of the non-zero second order magneto-transport effect for detection of Néel vector reversal.

Expanding $R(\theta)$ around the equilibrium angle θ_0 , we obtain:

$$V \approx \left[R(\theta_0) + \left(\frac{\partial R}{\partial \theta} \right)_{\theta_0} \Delta \theta \right] I, \quad (4.2)$$

where θ is the angle of the Néel vector and θ_0 its equilibrium position (i.e., at zero current). Assuming an alternating current $I = I_{AC} \sin(\omega t)$ and that the deviation from equilibrium is proportional to the current, $\Delta \theta = \delta_\theta I$, where δ_θ is a proportionality constant, we obtain:

$$V = R(\theta_0) I_{AC} \sin(\omega t) + \left(\frac{\partial R}{\partial \theta} \right)_{\theta_0} \delta_\theta I_{AC}^2 \sin^2(\omega t). \quad (4.3)$$

From this expression, three components can be identified: a DC term (V_{DC}), a first-harmonic term ($V_{1\omega}$), and a second-harmonic term ($V_{2\omega}$):

$$V_{DC} = \frac{1}{2} \left(\frac{\partial R}{\partial \theta} \right)_{\theta_0} \delta_\theta I_{AC}^2, \quad (4.4)$$

$$V_{1\omega} = R(\theta_0) I_{AC} \sin(\omega t), \quad (4.5)$$

$$V_{2\omega} = \frac{1}{2} I_{AC}^2 \left(\frac{\partial R}{\partial \theta} \right)_{\theta_0} \delta_\theta \sin \left(2\omega t - \frac{\pi}{2} \right). \quad (4.6)$$

The first-harmonic signal $V_{1\omega}$, oscillating at the same frequency as the applied AC current, corresponds to the linear response of the system, whereas the second-harmonic signal $V_{2\omega}$, oscillating at twice the frequency, represents the second-order (nonlinear) response. Higher-order magnetotransport measurements, such as the second-harmonic voltage detection described here, have been widely used to probe current-induced torques in FMs and heterostructures^{210–212}, as well as in globally inversion-symmetric (but locally non-centrosymmetric) AFs such as CuMnAs¹⁵.

Although the DC component V_{DC} contains information similar to that of the second-harmonic signal, its measurement is often complicated by spurious effects and artefacts that obscure the desired response. Nevertheless, the DC component will be further exploited in Chapter 6 to investigate SOT-driven antiferromagnetic resonance and self-oscillations via spin rectification effects, as the system is driven in resonance, the V_{DC} can be large enough and thus detectable.

In practice, resistance signals are often analyzed instead of voltage signals, defined as $R_{1\omega} = \frac{V_{1\omega}}{I_{AC}}$ and $R_{2\omega} = \frac{V_{2\omega}}{I_{AC}}$. The first-harmonic resistance $R_{1\omega}$ is independent of the probing current, whereas the second-harmonic resistance $R_{2\omega}$ varies linearly with it. This linear dependence of $R_{2\omega}$ on I_{AC} enables the identification and separation of the genuine signal from artefacts, such as those originating from imperfections in the AC source that generate higher-frequency components. Because $R_{1\omega}$ does not depend on the current, it can be directly compared across measurements at different probing currents, aside from minor variations caused by Joule heating. In contrast, second-harmonic resistance values should be normalized by the probing current to allow consistent comparison between measurements. Lastly, Thermoelectric effects arising from Joule heating are proportional to the square of the current and thus appear in the second-harmonic signal. Examples include the ANE and the SSE. Depending on the objective of the experiment, these effects may be regarded either as artefacts that obscure the desired signal or as useful mechanisms for signal detection.

4.2.3 Micromagnetic simulations of domain wall motion

Micromagnetic simulations were performed to compare DW dynamics in a single FM layer and a SAF, with particular focus on the influence of pinning centers

and stray fields on DW propagation. The simulations were carried out using Mumax3²¹³ over an area of $2\text{ }\mu\text{m} \times 0.5\text{ }\mu\text{m}$, discretized into $2048 \times 512 \times 3$ cells for the SAF and $2048 \times 512 \times 1$ cells for the FM layer. The SAF was modeled as two identical ferromagnetic layers sharing the same material parameters, with a saturation magnetization $M_s = 1.422 \times 10^6\text{ A/m}$ and a Gilbert damping $\alpha = 0.1$. The antiferromagnetic interlayer coupling and uniaxial anisotropy were adjusted to reproduce the experimental spin-flip fields along and perpendicular to the easy axis, resulting in an interlayer coupling of 0.2 T and a perpendicular anisotropy $K_u = 1.96 \times 10^6\text{ J/m}^3$. A SHA of 0.15 was assumed.

4.3 Results and discussion

4.3.1 Identification and quantification of magnetoresistances

In order to quantify the magnetoresistive effects in the system, first-harmonic (linear response) signals were analyzed during both magnetic field sweeps and field rotations, under a large saturating field of 2 T, well above the spin-flip field. As shown in [figure 4.4a](#), the out-of-plane field sweep reveals the linear responses of both transverse ($R_{xy}^{1\omega}$) and longitudinal ($R_{xx}^{1\omega}$) resistances. The longitudinal resistance exhibits a GMR-like stepwise behavior, characterized by a low-resistance state when the magnetic moments are aligned after the spin-flip transition, and a high-resistance state otherwise. In contrast, the transverse resistance is dominated by the AHE, where four distinct magnetic configurations can be identified.

Although the sample is nominally compensated, minor asymmetries between the ferromagnetic layers result in a nearly compensated SAF. These differences likely originate from structural variations in the growth conditions of the individual layers, leading to slight deviations in magnetic moment or anisotropy. Consequently, upon reducing the out-of-plane field and crossing the spin-flip threshold (approximately 0.15 T), the system stabilizes in a deterministic uniform antiparallel configuration. This unintentional asymmetry facilitates controlled preparation of magnetic states via external fields, enabling investigation of the device response in a single-domain regime. The two antiparallel configurations correspond to the states achieved after saturating the sample in positive and negative out-of-plane fields. Importantly, such slight layer inequalities do not

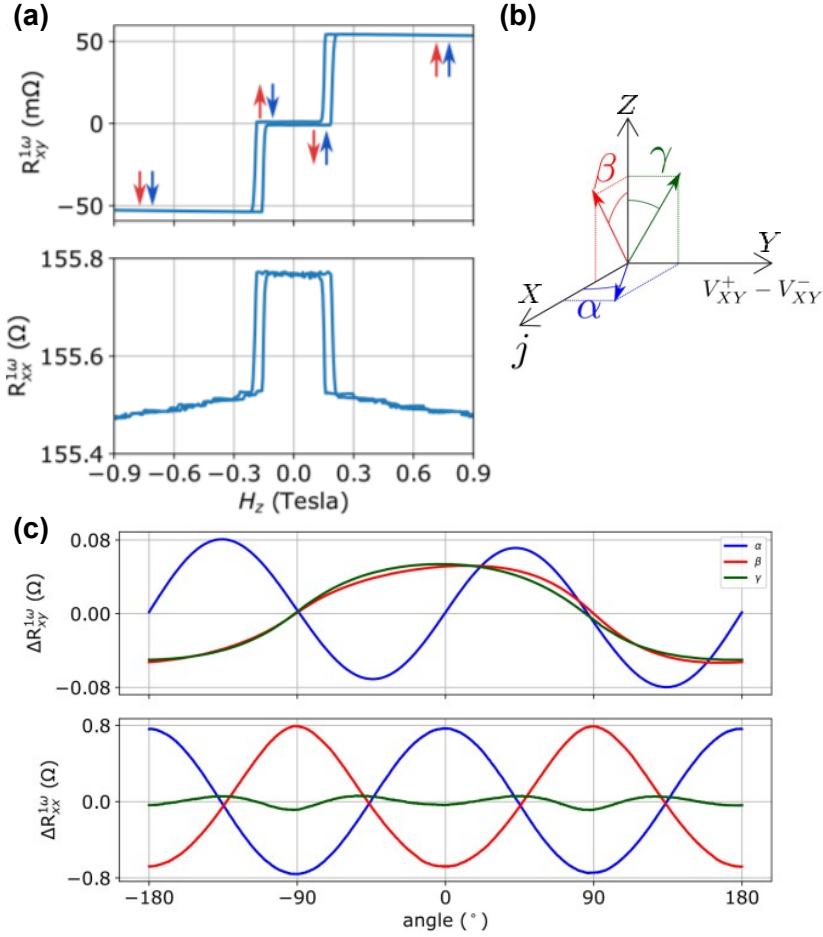


Figure 4.4: (a) Linear-response of both transversal ($R_{xy}^{1\omega}$) and longitudinal ($R_{xx}^{1\omega}$) resistances measured during an out-of-plane field sweep. No distinction between the two antiparallel states is observed, besides a small signal in $R_{xy}^{1\omega}$ due to an unintentional asymmetry between the ferromagnetic layers yielding a small net magnetic moment. (c) Linear responses $R_{xy}^{1\omega}$ and $R_{xx}^{1\omega}$ as a function of three primary angle rotations, depicted in (b).

compromise the memory functionality.

From the out-of-plane field sweeps, the contributions from AHE and GMR were extracted, while angular-dependent measurements were used to identify additional magnetoresistive effects. These angular scans were performed under a 2 T field to ensure a single-net magnetization state, effectively approximating the system to a single ferromagnetic layer. The resulting data were analyzed using a Fourier series decomposition,

$$f(x) = \sum_{n=0}^N a_n \sin(nx) + b_n \cos(nx), \quad (4.7)$$

allowing precise separation of the various contributions to the magnetoresistive signals. A summary of the observed magnetoresistances is presented in [table 4.1](#), specifying the corresponding rotational geometry and signal type ($R_{xx}^{1\omega}$ or $R_{xy}^{1\omega}$) from which each amplitude per sheet was derived. The field rotation data are shown in [figure 4.4c](#), with the angular conventions defined in [figure 4.4b](#).

Table 4.1: Summary of magnetotransport measurements and corresponding signal amplitudes.

Operation	Signal	MR Type	Amplitude per sheet
β -rotation	$R_{xx}^{1\omega}$	AMR	72 m Ω
γ -rotation	$R_{xx}^{1\omega}$	cAMR	5.6 m Ω
γ -rotation	$R_{xx}^{1\omega}$	SMR	1.1 m Ω
B_z sweep	$R_{xy}^{1\omega}$	AHE	54 m Ω
B_z sweep	$R_{xx}^{1\omega}$	GMR	25 m Ω (0.16%)

4.3.2 SOT-driven switching and detection of Néel order in SAF

To distinguish between the two anti-parallel (AP) states, we probed higher-order resistivity components that are independent of the net magnetic moment. An AC current of 5 mA (RMS) at a frequency of 163 Hz was applied through the Hall-bar device, and the longitudinal second-harmonic signal ($R_{xx}^{2\omega}$) was analyzed. As shown in [figure 4.5a](#), a clear contrast between the two AP states is observed in the $R_{xx}^{2\omega}$ response during magnetic field sweeps along the current direction (H_x).

To ensure unambiguous identification of the magnetic configuration, data were collected only during field sweeps from 0 to ± 2 T, thereby avoiding uncertainties in the magnetic order that can arise when crossing the in-plane spin-flip field (approximately 1 T) from high-field states, since both AP states are energetically equivalent. Before each of the four field sweeps, an out-of-plane field of ± 0.3 T was applied and subsequently removed to initialize the system in one of the two AP configurations.

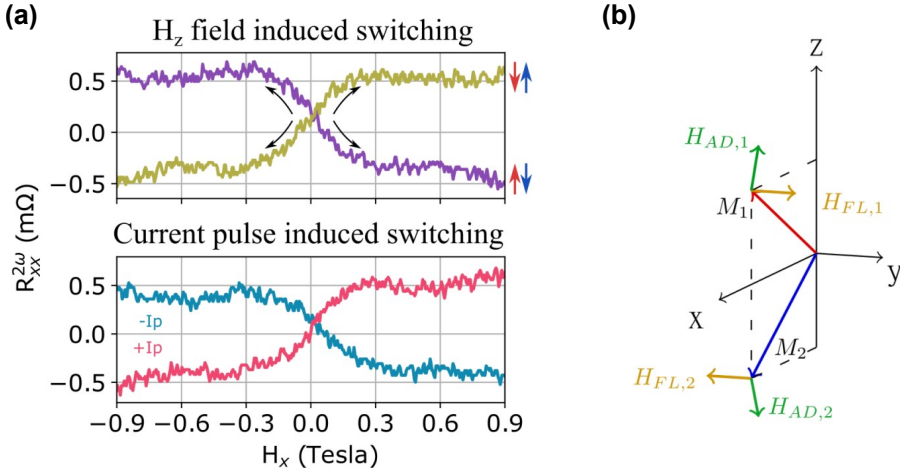


Figure 4.5: Second-order magnetoresistance in the out-of-plane SAF for detection of the Néel order. (a) Second-harmonic longitudinal signal ($R_{xx}^{2\omega}$) measured during in-plane field sweeps. To ensure a well-defined magnetic configuration, each sweep was initiated from zero field following the application of an out-of-plane field of ± 0.3 T, which fully magnetizes the sample into a specific antiparallel state. Comparable signals are obtained when applying an electrical pulse of 1.83×10^7 A/cm² and 10 ms long at $H_x = -0.9$ T instead of the out-of-plane presetting field, with the AP state determined by the pulse polarity. (b) Schematic illustration of the current-induced fields, highlighting the staggered nature of the antidamping field and its out-of-plane component, which energetically stabilizes one of the antiferromagnetic configurations.

From the expression of the second-harmonic signal (equation (4.6)), two primary conditions must be satisfied for it to be nonzero. First, the current-induced field must cause a deviation of the magnetic moments from their equilibrium position, i.e., $\delta_\theta \neq 0$. In other words, the current-induced field should not be collinear with the magnetic moments. Second, the resistance must vary

for positive and negative current directions, which requires $\left. \frac{\partial R}{\partial \theta} \right|_{\theta_0} \neq 0$. Consequently, in an out-of-plane magnetized system with in-plane current flow, no second-harmonic signal is expected at zero field, since this corresponds to a symmetric point (either a maximum or minimum) of the dominant magnetoresistive responses – AHE, GMR, and AMR. As shown in [figure 4.5a](#), at zero field the two magnetic states cannot be detected or distinguished, confirming this theoretical expectation.

Two types of current-induced effective fields can arise in the system. The first is the perpendicular, or FL, effective field $\mathbf{H}_{\text{FL}}$, which is oriented along the spin polarization $\boldsymbol{\sigma}$ and gives rise to the FL torque. The second is the DL effective field $\mathbf{H}_{\text{DL}}$, proportional to $\mathbf{M} \times \boldsymbol{\sigma}$. Because the probing current oscillates at a frequency much lower than the intrinsic magnetic precession frequencies, the magnetic moments can be considered to follow their instantaneous effective fields adiabatically. An illustration of these effective fields is shown in [figure 4.5b](#) for the case of an external magnetic field applied along the x direction. Here, $\mathbf{H}_{\text{FL}}$ is inherently staggered and tilts the sublattice moments within the plane, which, for an external field along x , produces no second-harmonic response. Thus, in this configuration, the FL contribution to the $R_{xx}^{2\omega}$ signal is negligible. In contrast, the DL field $\mathbf{H}_{\text{DL}}$ is non-staggered at zero field, pointing along $\pm x$. As the magnetic moments cant toward the x direction under an applied field, $\mathbf{H}_{\text{DL}}$ acquires a staggered out-of-plane (z) component. This component modulates the resistances (both AMR and GMR) differently for positive and negative current polarities, leading to the observed $R_{xx}^{2\omega}$ signal. Therefore, the second-harmonic longitudinal signal originates from the current-induced DL field. Detection of the two AP states becomes possible because $\mathbf{H}_{\text{DL}}$ reverses sign between the two antiferromagnetic configurations, tilting the sublattice moments in opposite directions for the same current polarity. Reversing the magnetic order thus produces signals of equal magnitude but opposite phase, as confirmed by the symmetric behavior observed when either the external field or the AP state is inverted.

Beyond detecting the two AP states, the staggered out-of-plane component of the current-induced effective field, $\mathbf{H}_{\text{DL}}$, can be exploited to achieve polarity-dependent electrical switching between these states. In [figure 4.5a](#), the electrical switching is compared with conventional out-of-plane field writing. Current pulses of 1.83×10^7 A/cm² and 10 ms duration were applied at -0.9 T. First, regardless of the initial state set by the external field, the pulse polarity uniquely

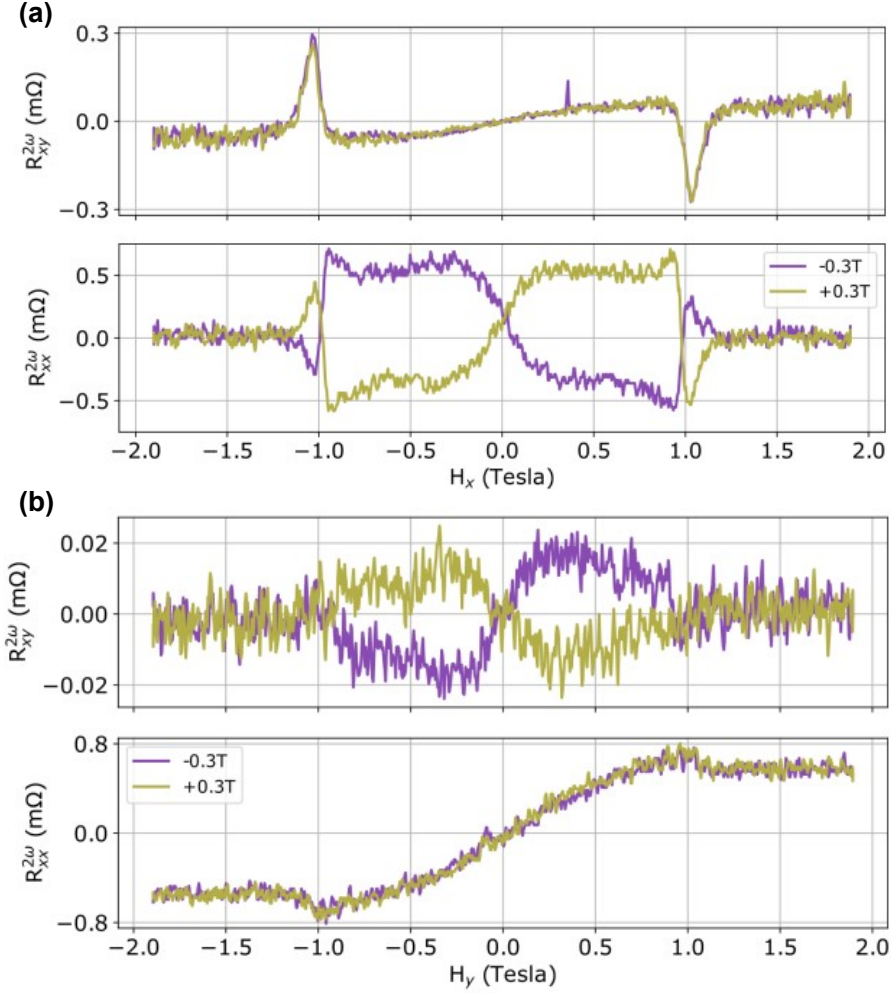


Figure 4.6: Nonlinear responses in SAF as a function of H_{extx} and H_y . The second harmonic signals, $R_{xy}^{2\omega}$ and $R_{xx}^{2\omega}$, were measured during field sweeps along the x and y directions, shown in (a) and (b), respectively. Prior to each sweep, a uniform antiparallel state was prepared by applying an out-of-plane field.

determines the final magnetic configuration. Second, the amplitude of the resulting signal is comparable to that obtained with field writing, indicating a complete magnetic reorientation across the volume between the longitudinal contacts.

This demonstrates that the system can operate as a magnetic memory, where both readout and writing are independent of any intentional imbalance between the ferromagnetic layers. Notably, this approach enables a two-terminal device geometry, in contrast to Hall-based measurement geometries. Furthermore, when the external field is applied along x , probing $R_{xx}^{2\omega}$ provides a signal free from thermal artifacts such as the ANE and SSE, which instead contribute to $R_{xy}^{2\omega}$.

In [figure 4.6](#), the second harmonic signals are shown for both field sweep directions. As before, each sweep was measured after applying an out-of-plane field to select one of the AP states. Thermal contributions are observed in $R_{xy}^{2\omega}$ and $R_{xx}^{2\omega}$ when the external field is applied along x and y , respectively. These background signals arise from the ANE and/or the SSE, which exhibit the same symmetry and are thus difficult to disentangle. They originate from an out-of-plane thermal gradient induced by Joule heating, with a total thermal contribution of approximately 0.055 mΩ per sheet. Considering that there are ten squares between the longitudinal contacts, the measured signal is enhanced tenfold.

It is noteworthy that detecting the AP state by applying an external field along the x -axis and measuring $R_{xx}^{2\omega}$ is not the only option. Alternatively, one can apply the field along the y -axis and measure $R_{xy}^{2\omega}$, as illustrated in [figure 4.6](#). However, this approach has several limitations. First, the signal magnitude is smaller and sensitive to the cross area of the device where the current density is non-uniform. More importantly, while the order parameter can still be detected, electrical switching cannot be achieved. From a macrospin perspective, this is due to the absence of a staggered component of the current-induced effective field along the easy axis.

When the field is applied along x , $R_{xy}^{2\omega}$ exhibits peak features near the spin-flip field. These peaks can be understood as a consequence of the flattening of the energy barrier²⁰³: with an increased in-plane field, the barrier between equilibrium states is reduced, allowing the current-induced field to induce a larger displacement or even switching between minima, resulting in a stronger nonlinear response. In the transverse contacts, the peak in $R_{xy}^{2\omega}$ arises from the AHE due to the small net moment. A similar effect is expected in the longitudinal

contacts, $R_{xx}^{2\omega}$, where the signal originates from the AMR. However, due to a slight misalignment of the in-plane field, a switching event is observed in the same field range where the peak is expected.

4.3.3 Realization of an analog multilevel antiferromagnetic memory in SAF

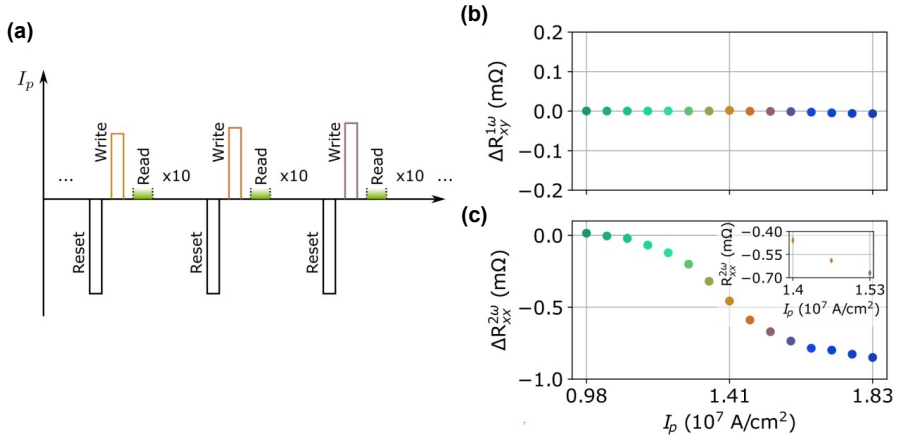


Figure 4.7: Multilevel electrical switching in the SAF. (a) Pulse sequence used to probe the dependence of the switching signal on the pulse amplitude, as well as its reproducibility. (b,c) Change in longitudinal resistance as a function of the pulse amplitude at $H_x = 0.9$ T for the first and second harmonics, $R_{xx}^{1\omega}$ and $R_{xx}^{2\omega}$, respectively. Each point represents the average over 10 pulses, and between writing pulses the magnetic state is reset by a negative pulse of -1.83×10^7 A/cm². Error bars correspond to the standard deviation of the mean value, which is highlighted in the inset.

The change in magnetoresistance as a function of the applied electrical pulse amplitude is presented in [figure 4.7](#). As previously discussed, the first-harmonic longitudinal signal, $R_{xx}^{1\omega}$, shows no discernible variation, reflecting its inability to distinguish between the two antiferromagnetic states ([figure 4.7b](#)). In contrast, the second-harmonic response, $R_{xx}^{2\omega}$, exhibits a clear dependence on the current density of the writing pulse ([figure 4.7c](#)). Notably, the response does not exhibit a step-like behavior characteristic of a complete reorientation of the magnetic order at a single critical current. Instead, a smooth, continuous curve is observed

between the two fully saturated states, with multiple intermediate levels. This behavior arises from partial switching across the sample, enabling coexistence of both AP states in varying proportions. Consequently, the system can function as a multilevel, or analog, memory device, with the number of distinguishable states ultimately limited by the measurement signal-to-noise ratio. In ideal AF films, the domain structure is dominated by magnetic anisotropy, and a single-domain ground state is expected due to negligible stray-field effects. However, in real devices, current-induced domain nucleation preferentially occurs at local defects, predominantly at the device boundaries in this case. Once formed, the inverted AF domains, created through SOT-driven domain nucleation and DW propagation, remain stable after the pulse is removed. The final size and position of these domains are determined by the minimization of DW length and pinning effects.

The writing procedure was performed under an external magnetic field along the Hall bar direction with $H_x = 0.9$ T, close to the hard-axis spin-flip field of approximately 1.1 T. Prior to each writing pulse, a strong negative-polarity pulse was applied to reset the device into a uniform AP state (inset of [figure 4.7c](#)). At each current density, the measurement was repeated ten times, and the mean value over 1 minute was plotted, with error bars corresponding to the standard deviation. The reproducibility of the switching process is confirmed by the narrow error bars, indicating only minimal variation in the nonlinear response between successive measurements.

4.3.4 Correlation between analog information and antiferromagnetic domains

To demonstrate that the nonlinear response signal, $R_{xx}^{2\omega}$, is directly related to the relative sizes of reversed antiferromagnetic domains, we compare the measured second harmonic signal with the domain distribution imaged using magnetic force microscopy (MFM). The nonlinear signal corresponding to a multidomain state prepared by a weak electrical pulse is shown in [figure 4.8a](#). Prior to this, two reference states with oppositely saturated Néel vectors were electrically written, and then a pulse of reduced magnitude generated a multidomain configuration. The resulting nonlinear response yields an analog signal of 0.68 ± 0.05 .

The longitudinal second harmonic signal originates from the region between

the two longitudinal contacts, which was scanned by MFM. [figure 4.8b](#) shows a section of the main channel of the Hall device, where two distinct regions, colored blue and green, correspond to domains with opposite Néel vector orientation. By integrating the areas of the two types of domains, we obtain a ratio of 0.75 ± 0.01 , in good agreement with the electrical measurement of $R_{xx}^{2\omega}$.

The multidomain state sensed by the nonlinear magnetoresistance is proportional to the difference between the areas of the reversed domains, A_{down} , and the total domain area $A_{\text{up}} + A_{\text{down}}$, such that $\frac{A_{\text{up}} - A_{\text{down}}}{A_{\text{up}} + A_{\text{down}}}$. Accordingly, the analog signal of the memory shown in [figure 4.7c](#) can be expressed as $\frac{A_{\text{down}}}{A_{\text{up}} + A_{\text{down}}}$. This demonstrates that the second harmonic response provides a quantitative readout of the multidomain configuration in the SAF.

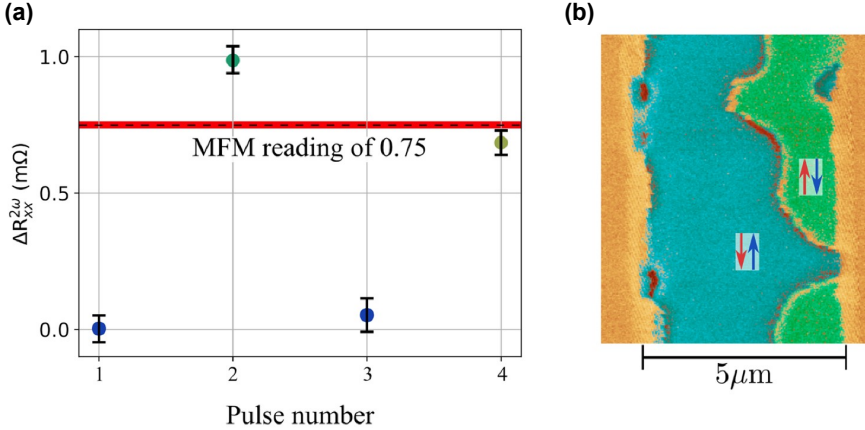


Figure 4.8: (a) Electrical preparation prior to the MFM scan. The first three pulses, of large magnitude and alternating polarity, drive the SAF into fully polarized single-domain states with opposite Néel vector orientations. The corresponding nonlinear magnetoresistance values serve as reference for the measurement performed after the fourth, weaker pulse, which generates a multidomain state. The expected ratio of domains with Néel vector up versus down, inferred from the electrical measurement, is $(68 \pm 5)\%$, while the analysis of the MFM scan yields a ratio of $(75 \pm 1)\%$. The dark dashed line indicates the ratio between integrated areas of oppositely oriented domains in the MFM image, and the red shaded region represents the error. (b) Section of the MFM scan highlighting large antiferromagnetic domains, shown in blue and green (colors edited for clarity).

4.3.5 Non-volatility of the antiferromagnetic memory

The non-volatile character of the antiferromagnetic DW memory is evident from the stability of the multidomain configuration, as confirmed by MFM measurements performed several days after the electrical preparation and initial detection. This non-volatility originates from the energetically degenerate antiferromagnetic states with reversed Néel vectors, separated by a high magnetic anisotropy barrier. Electrical writing requires an in-plane magnetic field to ensure deterministic switching, and this field additionally reduces the energy barrier between the two AP states, thereby lowering the onset current required for efficient switching.

The robustness of this non-volatile functionality is further demonstrated in [figure 4.9](#). A multidomain state was generated using a weaker electrical pulse and detected via the nonlinear magnetoresistance signal, $R_{xx}^{2\omega}$. Initially, the reference range corresponding to the two fully saturated antiferromagnetic states was established using a sequence of three saturating current pulses ($-I_p$, I_p , $-I_p$) with $I_p = 1.83 \times 10^7$ A/cm² at $H_x = 0.9$ T. Subsequently, a positive pulse of reduced magnitude (1.6×10^7 A/cm²) created a multidomain state. After measuring the corresponding nonlinear resistance, the in-plane magnetic field was switched off. After 60 minutes, the field was reapplied, the multidomain state remeasured, and finally, two saturating pulses were used to recover the reference states. Comparison of the nonlinear resistance values immediately after preparation and 60 minutes later shows only minor deviations within the experimental error, confirming the long-term stability of the multidomain configuration even in the absence of the in-plane magnetic field.

4.3.6 Antiferromagnetic domain wall profile and it's SOT-driven motion

We attribute the formation of the multidomain state underlying the multilevel memory behavior to current-induced domain nucleation at structural defects, followed by current-driven antiferromagnetic DW propagation. Both processes are governed by DL SOT. To demonstrate that the absence of stray magnetic fields accounts for the exceptional reproducibility in encoding analog information, micromagnetic simulations were carried out using the open-source, GPU-accelerated MuMax3 solver²¹³. The details of the micromagnetic simulations are provided in

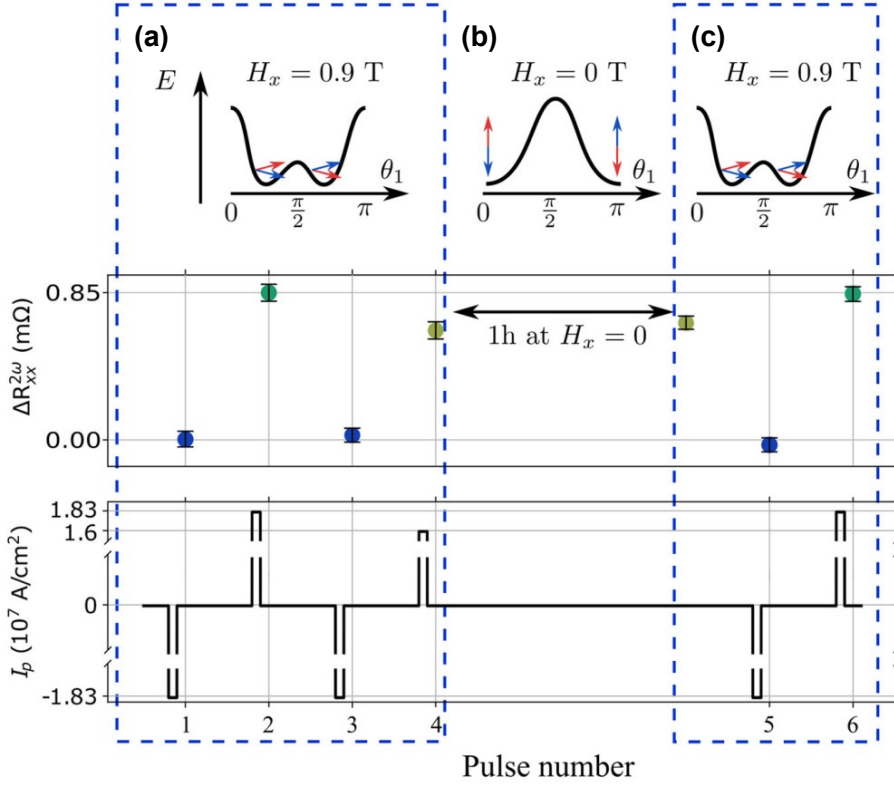


Figure 4.9: Non-volatile behavior of the multidomain state. (a) Electrical preparation of a multidomain state at $H_x = 0.9$ T. (b) The external magnetic field is switched off for 1 hour, and the multidomain state is subsequently remeasured after the field is turned on again. (c) Comparison of the nonlinear resistance signal, $R_{xx}^{2\omega}$, showing that the value measured after 1 hour coincides with the initial measurement within the error bars, with respect to the two fully saturated reference states. Above, a schematic illustration of the energy landscape as a function of the polar angles of both magnetic moments is shown for different external field magnitudes, assuming a cut at $\theta_2 = \pi - \theta_1$, where θ_i denotes the polar angle of the magnetization in layer i . Each point in the $R_{xx}^{2\omega}$ signal corresponds to the average of 1-minute readings, with error bars representing the standard deviation. Below, the sequence of electrical pulses applied during the measurement is illustrated.

???. For comparison, current-induced DW motion was investigated both in the SAF heterostructure and in a FM reference system composed of a single ferromagnetic layer adjacent to a nonmagnetic HM layer.

Figure 4.10a and figure 4.10b display the initial magnetic configurations consisting of two DWs within 500 nm wide and 2 μm long bars. Periodic boundary conditions were applied along the bar length to eliminate pinning effects arising from finite-size edges. In both the FM (figure 4.10a) and SAF (figure 4.10b) cases, the systems were initialized with a three-domain configuration and relaxed into low-energy states characterized by straight DWs of minimal length and approximately equal up/down domain populations. No artificial pinning or DW friction was introduced (e.g., through spatial variations in anisotropy), ensuring that the walls naturally relax into their energetically lowest configuration. The balanced distribution of domains in the FM system corresponds to the minimum total energy, accounting for the dipolar interaction associated with stray fields.

To introduce controlled pinning, regions of 20 nm diameter were defined, where the PMA was increased by 50% relative to the rest of the bar. These sites act as DW pinning centers by locally enhancing the reversal energy barrier. Analogous to the experimental procedure, an in-plane magnetic field H_x was applied along the bar axis and ramped up to 0.9 T. At the maximum field, an electrical current density of $j = 2 \times 10^7 \text{ A/cm}^2$ was applied, comparable to the experimental pulse amplitudes. The current pulse drives domain wall motion, which is halted at the pinning site. The pulse duration was chosen sufficiently long for the walls to reach a steady state before the current was switched off. The resulting magnetic configurations immediately after the pulse are presented in figure 4.10c and figure 4.10d for the FM and SAF systems, respectively.

In the FM bar, DW motion is strongly affected by both the pinning site and the long-range dipolar stray fields generated by the domains, leading to a complex stripe-like configuration. In contrast, for the SAF, wall propagation is impeded only locally by the pinning center, with no global distortion from stray fields. Additionally, the equal helicities of the twisted antiferromagnetic DWs within the film plane, where the DW normal $\mathbf{n}$ is perpendicular to the easy axis $\mathbf{e}_{\text{easy}}$ (detailed explanation is provided below in this section), prevent wall annihilation. As a result, adjacent walls approach each other until limited by the exchange energy cost of further compression. Figure 4.10e and figure 4.10f display the relaxed configurations after both the current and magnetic field have been turned off. In

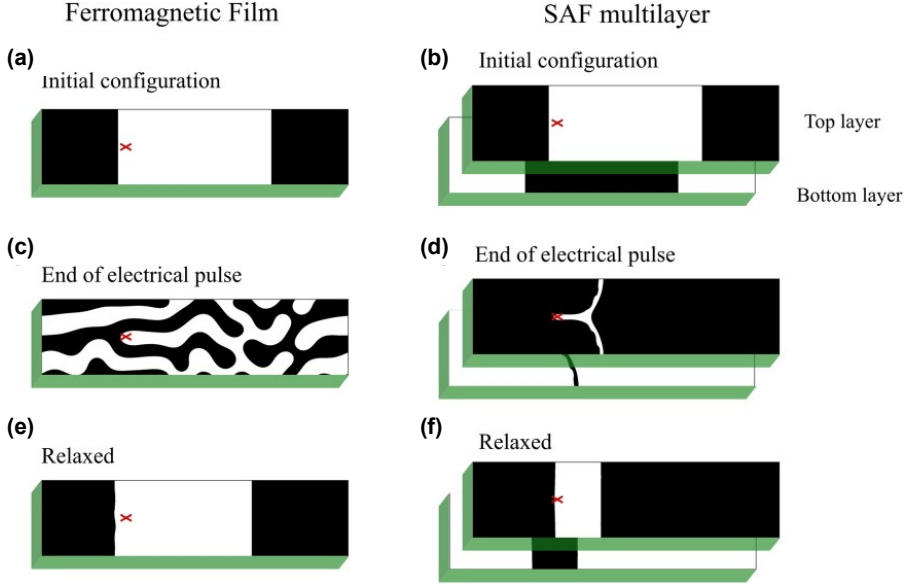


Figure 4.10: Comparison of current-induced DW motion in ferromagnetic and antiferromagnetic systems. The red cross marks the position of a pinning center, where the PMA is increased by 50% relative to the surrounding material. Black (white) regions denote areas with positive (negative) out-of-plane magnetization components. Panels (a) and (b) show the initial domain configurations in the ferromagnetic and antiferromagnetic systems, respectively. An in-plane magnetic field H_x is ramped along the bar direction up to 0.9 T, after which a current pulse of 20 ns duration is applied. The corresponding domain configurations at the end of the current pulse are presented in (c) and (d). Finally, the relaxed domain states after both the magnetic field and current are switched off are shown in (e) and (f). The simulations illustrate that in the FM, DW motion is strongly influenced by long-range stray fields, whereas in the AF, wall propagation and pinning are governed solely by local anisotropy variations.

both systems, straight DWs reappear as the DW energy is minimized. However, while the final FM configuration remains influenced by the long-range stray fields, the SAF domain pattern is governed solely by the local pinning site, where the wall stabilizes in its final position.

The present simulations include a single pinning center for conceptual clarity. In realistic systems with multiple pinning sites exhibiting a normal distribution of depinning potentials, the cumulative switching behavior follows an error-function dependence, assuming uncorrelated depinning events. Consistently, the experi-

mentally observed relationship between the fraction of reversed antiferromagnetic domains and the applied writing current density follows such an error-function trend (figure 4.7c).

The dark-light contrast observed at the antiferromagnetic domain boundaries in the magnetic MFM scans shown in figure 4.8b provides direct evidence of a localized dipolar magnetic field. This field is intrinsically small, spatially confined to the DW region, and oriented orthogonally to the domain boundary. The polarity of this dipolar field is uniquely determined by the relative Néel vector orientations of the two adjacent antiferromagnetic domains. Consequently, the dark-to-light contrast transition across the DW consistently proceeds from the domain appearing darker in the MFM image toward the lighter one. The physical origin of this contrast arises from the spatial displacement between the two antiferromagnetically coupled magnetic sublayers in the SAF structure. As schematically illustrated in figure 4.11a, a Bloch-like DW rotating in the z - y plane develops an additional magnetization component along the x -axis. This deviation twists the wall away from an ideal antiferromagnetic Bloch-wall configuration and induces a small, uncompensated magnetization component M_x . Furthermore, interfacial DMI at the Pt/Co-Ni interfaces, located at the top and bottom of the multilayer stack, may contribute additional chiral asymmetries. Depending on their relative sign and magnitude, these DMI contributions can either enhance or counteract the dipolar-field-induced twisting of the DW.

This interpretation is further supported by the micromagnetic simulations presented in figure 4.11. In figure 4.11b, the three components of the normalized magnetization vector are plotted along a horizontal line cut through the initial magnetic configuration shown in figure 4.11b. Both DWs in this configuration are analyzed, with the schematic in figure 4.11a corresponding to the wall centered at $x = 0.5 \mu\text{m}$. In addition to the expected $\pm M_y$ components associated with a Bloch-type wall, the simulation reveals a net negative M_x component. In contrast, the right-hand DW, located at $x = 1.5 \mu\text{m}$, exhibits a net dipole moment oriented along the $+x$ direction. Therefore, the local dipolar moments are directed toward the antiferromagnetic domain characterized by $\mathbf{M}_1 \uparrow$ and $\mathbf{M}_2 \downarrow$, denoted in the main text as the $\uparrow\downarrow$ state.

This analysis carries several implications. During the preparation of the multi-domain state shown in figure 4.8a, the application of a final current pulse with positive amplitude and a magnitude of $j = 1.59 \times 10^7 \text{ A/cm}^2$ drives the system

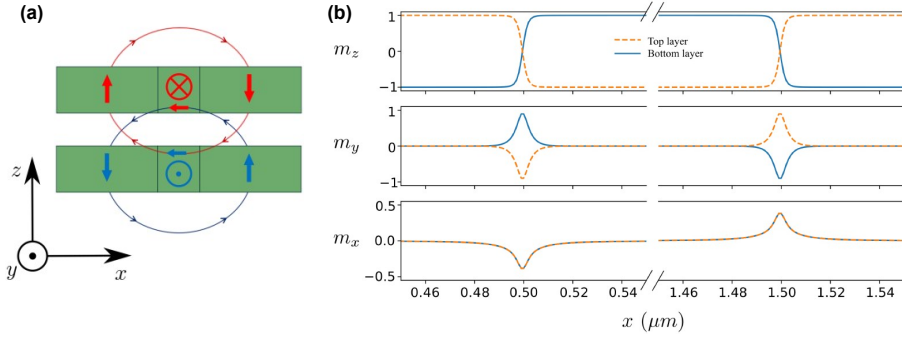


Figure 4.11: Twisted antiferromagnetic DW. (a) Schematic illustration of an antiferromagnetic DW in the SAF structure. Owing to the spatial displacement between the upper and lower antiferromagnetically coupled ferromagnetic layers, a local dipolar field emerges from adjacent domains. This dipolar interaction causes the otherwise Bloch-like DW profile to acquire an additional magnetization component perpendicular to the wall plane (along the x -axis). The resulting M_x components are parallel in both magnetic layers, thereby producing a net local dipole moment at the DW center. (b) Spatial variation of the normalized magnetization components (m_x , m_y , and m_z) along a horizontal line cut through the micromagnetic configuration presented in figure 4.10b. The simulation confirms the emergence of a finite M_x component, consistent with the twisted DW structure illustrated in panel (a).

toward a configuration where one domain type expands at the expense of its opposite counterpart. The MFM image reveals that the larger region of the device is predominantly occupied by the darker (or bluish) domain. Given that the orientation of the DW dipole moment is determined by the Néel vector configuration of adjacent domains, it follows that the darker (brighter) contrast corresponds to the antiferromagnetic state $\downarrow\uparrow$ ($\uparrow\downarrow$), respectively. Consequently, a positive current pulse applied under a positive in-plane magnetic field H_x results in the expansion of the $\downarrow\uparrow$ domain. The expanded domain type therefore depends sensitively on the combined polarity of both the applied current and H_x .

Notably, in figure 4.5a, current pulses are applied under a negative H_x , leading to the reversal of the expanded state relative to the previous configuration. A final point of comparison arises when contrasting the antiferromagnetic states written by current-induced switching with those written by out-of-plane external magnetic fields. Specifically, after applying a positive (negative) out-of-plane magnetic field exceeding the spin-flip field, the sample relaxes into a single-domain $\uparrow\downarrow$ ($\downarrow\uparrow$) configuration. This correspondence underscores the deterministic

control of antiferromagnetic domain states by the polarity of current and external field conditions.

4.4 Conclusion

In this work, we have demonstrated a non-volatile analog memory device operating at room temperature, based on a magnetically compensated SAF heterostructure. In this system, information is encoded in the ratio of reversed antiferromagnetic domain areas. The writing process relies on current-induced motion of antiferromagnetic DWs, while the readout is achieved via nonlinear magnetotransport measurements.

Biological synapses modify their synaptic weights according to the frequency and timing of input stimuli, a phenomenon known as spike time dependent plasticity (STDP). In contrast, our SAF-based antiferromagnetic analog memory exhibits a distinct form of plasticity *pulse-amplitude-dependent plasticity* in which the synaptic weight undergoes persistent strengthening or weakening depending on the amplitude of the applied current pulses. Modification of the stored analog state occurs only when the magnitude of a new pulse exceeds that of all previous pulses, or when the pulse polarity is reversed relative to the last applied pulse. We attribute the formation of the multidomain state to current-induced domain nucleation at structural defects, followed by DW propagation driven by DL SOT. The absence of stray magnetic fields in the SAF structure is a key factor enabling the remarkable reproducibility and stability of the analog information writing process.

Although the detection of reversed domain ratios via nonlinear magnetotransport produces relatively small signal amplitudes, an alternative and potentially more sensitive readout strategy can be realized using a three-terminal device geometry. In this configuration, one of the ferromagnetic layers in the SAF forms part of a MTJ, while the counter electrode serves as a magnetic reference layer. The resulting TMR signal directly reflects the ratio of reversed domains within the SAF, thus providing a quantitative measure of the stored analog information. Integrating such an SAF-based antiferromagnetic memory into an MTJ device could enable its use as a magnetic synapse interfaced with spintronic functional units such as STT oscillators or probabilistic p-bit devices, where the synaptic

weight would control the degree of spin polarization driving the oscillatory or stochastic behavior.

A further refinement of this concept involves replacing the ferromagnetic reference layer with a reference SAF structure. This reference SAF remains in a fully polarized single-domain state and is magnetically decoupled from the “free” SAF below. The resulting all-antiferromagnetic analog memory would feature a large-magnitude, linear-response TMR readout, while still being written by SOT-driven DW motion. Importantly, DL SOT-induced antiferromagnetic DW motion remains viable even when spin accumulation arises only at one interface of the SAF stack.

A system composed of two magnetically compensated SAFs exhibits zero or nearly zero net magnetization, preserving the advantages of a fully compensated spin structure. Such configurations are inherently robust against external magnetic perturbations and self-generated stray fields, offering enhanced stability and scalability for neuromorphic and probabilistic computing applications. Finally, it is worth emphasizing that the principles demonstrated here are not limited to SAF systems. Analog antiferromagnetic memory functionality could also be realized in real AFs or in recently discovered *altermagnets*⁹⁷, where current-induced DW motion is likewise anticipated. These results pave the way toward a new class of fully compensated, SOT-driven analog spintronic devices that combine non-volatility, analog tunability, and field insensitivity – key ingredients for future brain-inspired and probabilistic information technologies.



Spin-orbit enabled tr-MOKE imaging of antiferromagnetic domains

As shown in [chapter IV](#), SOT control of antiferromagnetic textures offers strong potential for racetrack memory and neuromorphic computing. In this chapter, we report the direct imaging of current-induced antiferromagnetic domain dynamics in a constriction-based in-plane SAF using stroboscopic tr-MOKE microscopy. Although polar MOKE is usually insensitive to in-plane SAF order, a staggered DL spin-orbit field acting oppositely on the two magnetic layers generates a dynamic out-of-plane moment, enabling real-space detection of antiferromagnetic domains. Phase-resolved measurements reveal that the Néel order in adjacent domains oscillates 180° out of phase, producing oppositely signed tr-MOKE signals. The domain and DW orientation is set by a pre-applied magnetic field, while DC current injection enables reversible DW motion observed dynamically. Finally, a localized excitation mode at the constriction shows how geometry confines antiferromagnetic dynamics.

Keywords: tr-MOKE microscopy, Néel-order oscillations, domain-wall motion, SOTs.

This chapter is based on: **Rout, P.K.**, Vilsmeier, F. et al. Imaging of antiferromagnetic domain dynamics enabled by staggered spin-orbit fields. Under preparation.

5.1 Introduction

AFs were historically overlooked due to detection limitations, but recent advances highlight their ultrafast terahertz dynamics and vanishing stray fields, driving the emergence of antiferromagnetic spintronics^{3,5,31}, as outlined in [chapter I](#). As discussed in [chapter IV](#), current-induced SOT enables efficient magnetization switching and current-driven DW motion in SAFs with PMA, where the antiferromagnetic states can be electrically detected via nonlinear magnetotransport²⁰³. More generally, current-induced SOT^{124,125} has emerged as a central mechanism in spintronics enabling efficient magnetization switching^{15,17,22,80,214}, DW and skyrmion motion^{215–218}, magnetization self-oscillations^{79,80}, and spin-wave excitations^{219,220}. In the $\mathcal{PT}$ -symmetric antiferromagnet CuMnAs, Néel SOT enables reversible DW motion, while antiferromagnetic domain structures can be imaged using X-ray photoemission electron microscopy (XPEEM) with contrast arising from X-ray magnetic linear dichroism (XMLD)¹⁸. Current-induced DW motion with velocities up to 750 m/s has been demonstrated in a SAF with PMA, as measured using polar magneto-optical Kerr effect (p-MOKE)³². In such systems, the MOKE contrast is typically weak and arises from the net uncompensated magnetization of the SAF structure. Imaging antiferromagnetic domains presents a longstanding experimental challenge, even for the static case, and their time-resolved, dynamic real-space visualization is still lacking. Fully compensated AFs provide little magneto-optical contrast, and in-plane systems are especially elusive: the widely used p-MOKE geometry is intrinsically insensitive to in-plane FM or AF order, while longitudinal MOKE requires scanning-based implementations that are experimentally demanding and often lack sufficient sensitivity.

Very recently, a parallel study demonstrated that in an α -Fe₂O₃/Pt bilayer, a current-modulated MOKE technique can access SOT effects in an in-plane AF. In this system, the damping-like SOT effective field tilts the small canted weak ferromagnetic moment of α -Fe₂O₃, generating an out-of-plane component that becomes detectable in the polar MOKE geometry. In this case, the p-MOKE resolves the dynamics of the weak ferromagnetic moment and allows quantitative determination of the SOT²²¹.

In this study, we directly image current-induced antiferromagnetic domain dynamics in a constriction-based, in-plane easy-plane SAF using stroboscopic

tr-MOKE microscopy. The SAF serves as a simplified yet tunable model platform for exploring fundamental aspects of antiferromagnetic spin dynamics. Although p-MOKE is intrinsically insensitive to in-plane SAF order, a staggered antidamping spin-orbit field acting with opposite sign on the two magnetic layers generates a dynamic out-of-plane component, enabling real-space optical detection of antiferromagnetic domains. The tr-MOKE signal reverses sign between adjacent domains with opposite Néel vector orientations. Phase-resolved measurements reveal that this contrast originates from a 180° phase shift in the Néel vector oscillations across domain boundaries, consistent with the symmetry of the DL effective field. The domain, and consequently DW, orientation is set by the direction of a prior saturating magnetic field. We attribute this behavior to spin-flop-induced reorientation of the SAF magnetization and to the tendency of DWs to minimize magnetic charge. Above a critical DC current threshold, we observe reversible current-driven DW motion. In addition to the non-resonant excitation mentioned above, a pronounced coherent excitation mode emerges that is strongly confined within the micron-scale constriction when the magnetic field is applied collinearly with the current-carrying bar.

5.2 Experimental Methods

5.2.1 Sample and device fabrication

For these studies, we prepared SAFs consisting of two Py ($Ni_{81}Fe_{19}$) ferromagnetic layers exhibiting easy-plane anisotropy. A schematic of the SAF multilayer stack is shown in [figure 5.1a](#). The antiferromagnetic coupling between the two Py layers was achieved via the RKKY interaction, through a 0.34 nm Ir spacer layer. A 5-nm Ta seedlayer was included underneath the stack to improve adhesion. The whole Ta(5 nm)/Py(5 nm)/Ir(0.34 nm)/Py(5 nm)/Pt(5 nm) multilayer was deposited at room temperature by d.c. magnetron sputtering in an ultrahigh-vacuum BESTEC system. Argon was supplied at a flow of 20 standard cubic centimeters per minute (sccm), resulting in an Ar atmosphere at 0.4 Pa (3 mTorr). The substrates were Si wafers with a 100-nm thermally oxidized SiO_2 layer. During deposition, the substrate was rotated at ≈ 60 rpm. The d.c. sputter power was set to 50 W for Py, Pt, and Ta, and 40 W for Ir corresponding to the growth rates of 0.255 nm/s for



Py, 0.1 nm/s for Pt, 0.068 nm/s for Ta, and 0.056 nm/s for Ir. Layer thicknesses were controlled via deposition time, and sputter rates were calibrated in advance by X-ray reflectivity.

The constrictions of the devices were defined by electron-beam lithography (EBL) using a positive e-beam resist (AR-P 6200). The exposed regions were etched by argon ion milling in a reactive-ion etching–inductively coupled plasma (RIE–ICP) system. Electrical contacts were fabricated using a standard lift-off process following thermal evaporation of a Cr/Au bilayer (5 nm/80 nm). Figure 5.1b shows the optical micrograph of a patterned bar shaped device containing a single notch-shaped microconstriction to enhance the current density and spatially confine magnetic excitations.

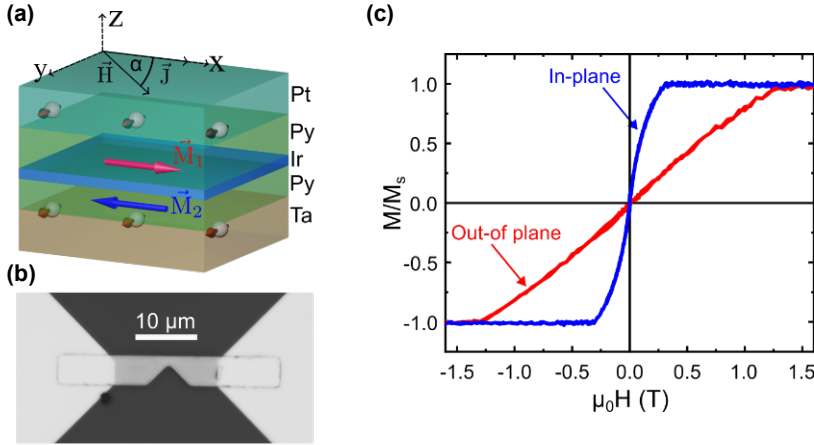


Figure 5.1: (a) Schematic of the SAF multilayer stack comprising two permalloy (Py: $\text{Ni}_{0.81}\text{Fe}_{0.19}$) layers antiferromagnetically coupled via a 0.34-nm Ir spacer, sandwiched between Pt and Ta layers. The green arrows indicate the spin polarization directions injected from the heavy-metal interfaces when a charge current flows through the multilayer stack. Red and blue arrows represent the magnetization vectors ($\mathbf{M}_1$) and ($\mathbf{M}_2$) of the upper and lower Py layers, respectively. The in-plane magnetic field ($\mathbf{H}$) and charge current ($\mathbf{J}$) are applied at an angle α . (b) Optical micrograph of the fabricated constriction device, displaying a notch-shaped central constriction with a width of 1 μm . (c) Magnetization curve of the SAF: Normalized magnetization as a function of the magnetic field H for in-plane (blue) and out of plane (red) directions measured in an unpatterned SAF.

Magnetometry measurements were performed using a commercial Microsense EZ7 vibrating sample magnetometer equipped with an electromagnet providing

fields up to 1.8 T. Hysteresis loops were recorded with fields applied both in-plane and out-of-plane to probe the magnetic response of the SAF (figure 5.1c). The measurements reveal a pronounced hard-axis anisotropy oriented perpendicular to the film plane, with an anisotropy field of ~ 1 T. The in-plane hysteresis loop exhibits a distinct spin-flip transition at 0.3 T, marking the critical field at which the magnetizations of the two ferromagnetic Py layers align parallel to the applied field. This behavior is characteristic of strong RKKY exchange coupling mediated by the ultrathin Ir spacer. Importantly, the negligible hysteresis at zero field demonstrates excellent magnetic compensation between the two Py layers, whose nearly identical moments cancel, resulting in an almost vanishing net magnetization.

5.2.2 Scanning time-resolved magneto-optic Kerr effect microscopy

MOKE relies on magnetization-induced modifications of the polarization state of light reflected from a magnetic material. When linearly polarized light (a superposition of left- and right-handed circularly polarized light) is incident on a magnetic film, spin-orbit coupling and exchange-induced band splitting lead to different optical responses for left- and right-circularly polarized light^{7,8,III,222,223}. As a result, the reflected beam exhibits a rotation of its polarization plane (Kerr rotation) and an induced ellipticity. Depending on the relative orientation of light incidence and magnetization, three MOKE geometries are distinguished^{7,III}:

- **Polar MOKE:** sensitive to out-of-plane magnetization components. It is characterized by a change of rotation of polarization and an ellipticity of the reflected beam.
- **Longitudinal MOKE:** sensitive to in-plane magnetization parallel to the plane of incidence. Here, a rotation of polarization and a slight ellipticity of the reflected beam are observed.
- **Transverse MOKE:** sensitive to in-plane magnetization perpendicular to the plane of incidence, detected as an intensity change.

Macroscopically, the magneto-optical response is described by the permittivity tensor (ϵ)²²³. For a film with out-of-plane magnetization, the permittivity tensor



takes the form

$$\boldsymbol{\varepsilon} = \varepsilon_0 \varepsilon_r \begin{pmatrix} 1 & iQ & 0 \\ -iQ & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (5.1)$$

where ε_0 and ε_r denote the vacuum and relative permittivity, respectively, and Q is the complex magneto-optical parameter. The off-diagonal elements originate from spin-orbit interaction and are responsible for the Kerr effect. They lead to different complex refractive indices for right- and left-handed circularly polarized light^{7,8},

$$n_{\pm} \approx \sqrt{\varepsilon_r} \left(1 \pm \frac{Q}{2} \right). \quad (5.2)$$

Because these left- and right-handed circularly polarized light propagates with different phase velocities and attenuation constants, their superposition, corresponding to linearly polarized light, acquires a rotation and ellipticity upon reflection. The Kerr signal changes sign upon magnetization reversal, reflecting its odd in $\mathbf{M}$ symmetry. Consequently, conventional p-MOKE is insensitive to fully compensated collinear antiferromagnets, whose net magnetization vanishes and also to the in-plane FMs or AFs.

In this work, the probe beam is incident normal to the film surface, ensuring exclusive access to the p-MOKE configuration. To probe the SOT driven dynamic magnetization tr-MOKE technique is employed. Figure 5.2 illustrates the principle of the tr-MOKE experimental setup used in this thesis. tr-MOKE microscopy is a powerful optical pump-probe technique widely employed to investigate magnetization dynamics such as spin-wave excitations, vortex oscillations and spin Hall nano-oscillators^{224–231}. In the following, a brief overview of the tr-MOKE microscopy employed throughout this thesis is presented. The experiments were carried out in the group of Prof. C. H. Back at the Technical University of Munich. A detailed description of this specific tr-MOKE implementation can be found elsewhere^{229–231}.

In our experiments, magnetization dynamics is driven by SOTs from microwave excitation at frequency f , while the dynamic response is probed stroboscopically using ultrashort laser pulses. A mode-locked Coherent Mira 900 laser serves as the light source, delivering femtosecond pulses at a repetition rate $f_{\text{rep}} = 80$ MHz and a central wavelength $\lambda = 800$ nm. Since the pulse duration is much shorter

than one excitation period, each laser pulse samples a well-defined phase of the magnetization dynamics. By maintaining a fixed phase (relative delay) relation, δ , between the microwave excitation and the laser pulse train, a single phase of the oscillatory magnetization is detected. This requires the excitation frequency to be an integer multiple of the repetition rate i.e. $f = \nu f_{\text{rep}}$, $\nu \in \mathbb{N}$, ensuring phase-locked synchronization between excitation and probe.

Synchronization is achieved using a phase-locking module, *Synchrolock*, that stabilizes the laser repetition rate via an internal feedback loop. A fast photodiode monitors the pulse train, and a piezo-controlled cavity mirror compensates for drifts. The module provides a 10 MHz reference signal to the microwave generator, guaranteeing a constant phase relation between optical probe and radio-frequency (rf) excitation.

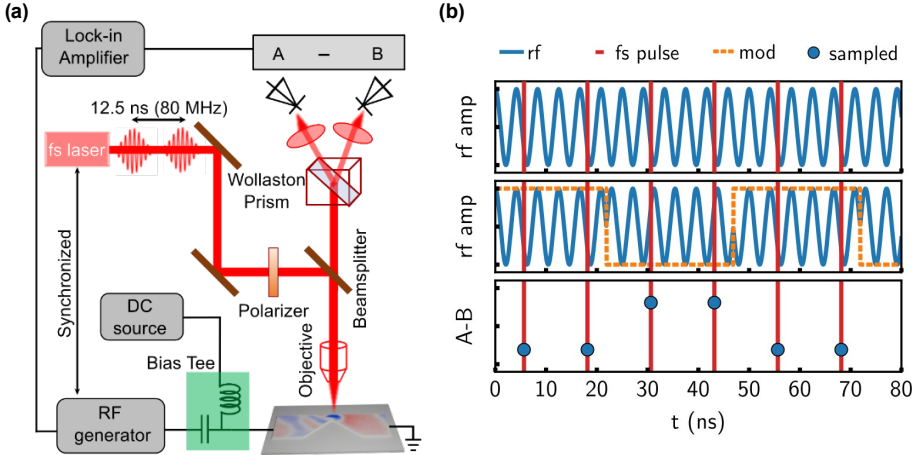


Figure 5.2: (a) Schematic illustration of the scanning time-resolved magneto-optical Kerr effect (tr-MOKE) experimental setup used to probe SOT driven magnetization dynamics. (b) Top: Illustration of the synchronization between the rf excitation (blue) and the probing laser pulses (red vertical lines). Middle: The kHz reference modulation (orange dashed line) and the resulting 180° phase-modulated rf signal. Bottom: The discrete magnetization values sampled stroboscopically by the laser pulses (blue dots), which are subsequently demodulated and detected by the lock-in amplifier. Adapted from [231].

The linearly polarized laser beam is focused onto the sample using an objective lens with numerical aperture $\text{NA} = 0.7$. The diffraction-limited spatial resolution is given by $d_{\text{res}} = \frac{\lambda}{2\text{NA}}$, which yields $d_{\text{res}} \approx 571 \text{ nm}$ for $\lambda = 800 \text{ nm}$. The reflec-

ted beam experiences a Kerr rotation proportional to the dynamic out-of-plane component of the magnetization when operating in the p-MOKE geometry. The polarization change is analyzed using a Wollaston prism, which separates the reflected light into two orthogonally polarized components detected by a balanced photodetector. The difference signal ($A - B$) between the two photodiode currents provides the Kerr signal and is directly proportional to the dynamic magnetization component. The sum signal ($A + B$) corresponds to the sample reflectivity and is used for normalization and detection of device edges during scanning. To enhance sensitivity, lock-in detection is employed. The rf excitation is phase-modulated by 180° at a low (kHz) reference frequency using frequency mixing. After amplification, the modulated signal is delivered to the wire-bonded SAF device mounted on a custom built high-frequency PCB. The excitation is delivered through the high-frequency port of a bias-tee, enabling simultaneous injection of rf and DC signals. The Kerr signal is demodulated at the reference frequency using a lock-in amplifier, whose in-phase (X) output is proportional to the dynamic out-of-plane magnetization component.

It is important to note that the measured Kerr amplitude corresponds to the magnetization at a fixed relative phase delay (δ) between excitation and probe. To reconstruct the full precessional amplitude and phase, an additional phase resolved measurement can be performed where the relative delay (δ) between the microwave current and the laser repetition rate can be varied from 0° to 360° at a fixed frequency to facilitate phase-sensitive detection of the dynamic Kerr signals.

For spatial imaging, the sample is mounted on a high-precision xyz piezoelectric stage, enabling two-dimensional scans of the laser spot across the surface. This allows mapping of the dynamic Kerr response with sub-micron spatial resolution. An in-plane magnetic bias field can be applied using a rotatable electromagnet.

5.3 Results and discussion

5.3.1 spin-orbit enabled polar MOKE detection of in-plane SAF domains

To access antiferromagnetic domains in the in-plane SAF using the otherwise insensitive p-MOKE geometry, we dynamically excited the SAF magnetization via SOT. A microwave current was applied through the device, and the resulting time-resolved Kerr response of the dynamic out-of-plane component of the net magnetization, $\Delta m_z(t)$, was measured in polar geometry. The dynamic Kerr signal amplitude can be expressed as $\tilde{\Theta}_K = \Theta_K e^{i\varphi}$, where Θ_K is the magnitude and φ is the phase of the complex dynamic Kerr response. As described in [section 5.2.1](#), the SAF multilayer consists of two 5 nm thick permalloy (Py: $\text{Ni}_{0.81}\text{Fe}_{0.19}$) layers antiferromagnetically coupled via a 0.34 nm Ir spacer through the RKKY interaction. The system exhibits nearly fully compensated behavior with in-plane easy-plane anisotropy, such that both magnetizations lie within the film plane. Consequently, no static p-MOKE signal is expected.

The SAF stack is sandwiched between a Ta seed layer and a Pt capping layer, which possess opposite SHAs. Upon applying a charge current, spin currents are generated at both heavy-metal interfaces via the spin Hall and Rashba-Edelstein effects (see [chapter II](#)) and are injected into the SAF with additive spin polarization. A microwave current produces an oscillating spin current at the same frequency, generating alternating FL ($\mathbf{H}_{\text{FL}}$) and DL ($\mathbf{H}_{\text{DL}}$) effective fields that drive magnetization dynamics. The resulting dynamic out-of-plane component m_z becomes detectable via polar tr-MOKE. As will be discussed below, it is specifically the staggered nature of the DL spin-orbit field $\mathbf{H}_{\text{DL}}$ that, in principle, enables the optical detection of the antiferromagnetic domains.

[Figure 5.3a](#) shows the tr-MOKE signal measured at a microwave excitation frequency of 4 GHz with a relative delay $\delta = 90^\circ$ between the microwave excitation and the laser pulse train, under zero applied magnetic field. Prior to measurement, a saturating magnetic field was applied collinearly with the current-carrying bar ($\alpha = 0^\circ$, where α denotes the azimuthal angle with respect to the bar direction, see [figure 5.1a](#)). The tr-MOKE image exhibits alternating red and blue regions, corresponding to opposite signs of the dynamic out-of-plane magnetization com-



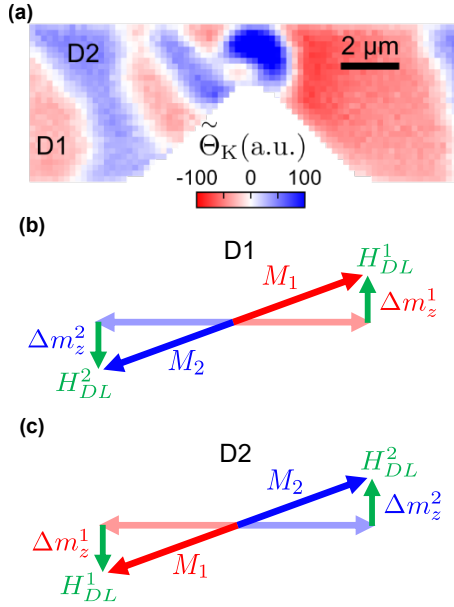


Figure 5.3: (a) tr-MOKE image showing the dynamic Kerr signal amplitude ($\tilde{\Theta}_K$) acquired at the microwave excitation frequency of 4 GHz and zero applied magnetic field. The relative delay (δ) between the laser repetition rate and the microwave excitation signal is 90° . D1 and D2 correspond to the two antiferromagnetic domains with reversed Néel order. (b),(c) Schematic illustration of the staggered DL spin-orbit fields (green arrows) for two antiferromagnetic domains, D1 and D2. The red and blue arrows represent the magnetizations of the two SAF layers. Upon reversal of the Néel vector, the staggered DL spin-orbit fields ($\mathbf{H}_{DL}^1$ and $\mathbf{H}_{DL}^2$) acting on each magnetic layer reverse sign accordingly. In an ideal fully compensated SAF, no net dynamic out-of-plane magnetization δm_z would arise, resulting in a vanishing p-MOKE response. However, asymmetries between the Pt/Py and Ta/Py interfaces can lead to an imbalance in the SOTs as well as MOKE response, giving rise to a finite net δm_z and thus a detectable polar Kerr signal, which reverses sign between the two oppositely oriented domains.

ponent, separated by narrow white lines with near-zero signal. We attribute the red and blue areas to antiferromagnetic domains D1 and D2 respectively, undergoing non-resonant oscillatory dynamics driven by the staggered DL spin-orbit fields generated by the microwave current. The white lines are interpreted as DWs separating adjacent domains.

The damping-like spin-orbit field induces a dynamic out-of-plane net magnetization whose sign depends on the orientation of the Néel vector (figures 5.3b

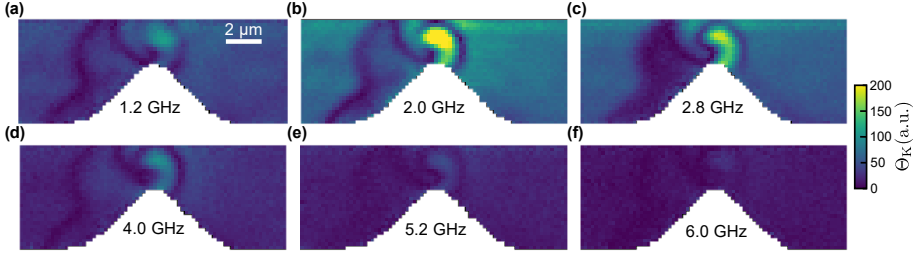


Figure 5.4: (a-f), Spatial maps of the tr-MOKE signal magnitude (Θ_K) at six excitation frequencies: 1.2, 2.0, 2.8, 4.0, 5.2, and 6.0 GHz. Prior to scanning, a saturating magnetic field was applied along the bar axis ($\alpha = 0^\circ$), followed by field removal. The 2.0 GHz map exhibits the strongest signal amplitude, indicative of resonance-like behavior.

and 5.3c). When the Néel vector is reversed between the two antiferromagnetic domain states, the staggered DL spin-orbit fields, $\mathbf{H}_{DL}^1$ and $\mathbf{H}_{DL}^2$, acting on the respective magnetic layers likewise change sign. Consequently, domains with opposite Néel vector orientations generate Kerr signals that are phase-shifted by 180° , enabling direct optical discrimination of antiferromagnetic domains separated by 180° domain walls.

In an ideal fully compensated SAF, no net dynamic out-of-plane magnetization δm_z would emerge, leading to a vanishing p-MOKE response. However, structural asymmetry between the Pt/Py and Ta/Py interfaces can produce an imbalance in both the SOT strength and the magneto-optical response, resulting in a finite net δm_z and thus a detectable Kerr signal. Notably, it has been reported that the Pt/Py interface exhibits a stronger MOKE response than the Ta/Py interface²³². In addition, a small uncompensated magnetization may further contribute to a finite polar MOKE signal, which reverses sign between oppositely oriented domains. However, the absence of any hysteresis at zero field from the magnetometry measurements (see figure 5.1c) excludes the MOKE response from the uncompensated moment.

The interpretation presented above is based on a non-resonant excitation of the antiferromagnetic domains driven by the staggered damping-like spin-orbit fields. However, the observed opposite red and blue tr-MOKE contrast could alternatively be attributed to a resonant excitation of a chiral antiferromagnetic eigenmode. At the present stage, this scenario cannot be conclusively established and requires further investigation. The emergence of such a chiral eigenmode would rely on

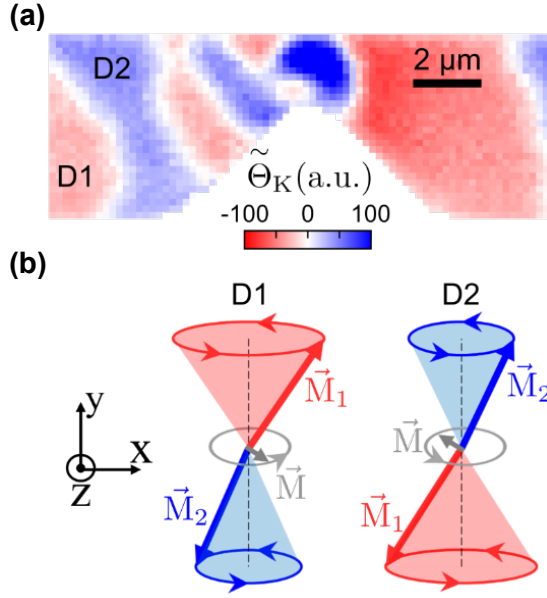


Figure 5.5: (a) tr-MOKE image showing the dynamic Kerr signal amplitude ($\tilde{\Theta}_K$) acquired at the microwave excitation frequency of 4 GHz and zero applied magnetic field. The relative delay (δ) between the laser repetition rate and the microwave excitation signal is 90° . D1 and D2 corresponds to the two antiferromagnetic domains with reversed Néel order. (b) Circularly polarized easy-axis antiferromagnetic eigenmode (right-handed) corresponding to the two antiferromagnetic domains D1 and D2 respectively. The reversal of the Néel order between D1 and D2 results in a 180° phase shift of the oscillations of net magnetization $\vec{M}$ for the right-handed eigenmode.

the presence of a weak in-plane uniaxial easy-axis anisotropy superimposed on the dominant in-plane easy-plane anisotropy. Such an additional anisotropy contribution may arise from device patterning. As discussed in [chapter III](#), an easy-axis AF below the spin-flop field supports circularly polarized left- and right-handed chiral eigenmodes. In the following, we outline how excitation of such a chiral mode of given chirality could also give rise to opposite polar tr-MOKE signals in adjacent antiferromagnetic domains.

The linearly polarized acoustic mode of a pure easy-plane AF vanishes at zero applied magnetic field, while the optical mode persists at much higher frequencies, well beyond our 4 GHz excitation. However, the frequency dependence of the tr-MOKE amplitude Θ_K measured at zero field ([figure 5.4](#)) reveals a pronounced

maximum at 2.0 GHz, with the signal nearly vanishing at 1.2 GHz and above 5.2 GHz. This resonance-like behavior suggests the presence of a weak in-plane easy-axis anisotropy superimposed on the dominant easy-plane anisotropy, possibly induced by device patterning. Such a uniaxial contribution can stabilize circularly polarized left- and right-handed eigenmodes even at zero magnetic field. In an ideal compensated AF, these two chiral modes are degenerate. In practice, small asymmetries between the two ferromagnetic layers of the SAF or a minute residual magnetic field can lift this degeneracy and select an eigenmode with fixed handedness. At 4 GHz, slightly detuned from the 2.0 GHz resonance, the tr-MOKE measurements may reveal coherent oscillatory dynamics consistent with excitation of such a chiral mode. It is important to note that, below the spin-flop field, the chiral modes are predominantly governed by the magnetic anisotropy fields, which are staggered between the two sublattice magnetizations. Owing to its staggered character, in our SAF system, the DL spin-orbit field is therefore ideally suited to excite these chiral eigenmodes.

In polar geometry, we detect the dynamic out-of-plane component of the net magnetization. The opposite red and blue contrasts observed in [figure 5.5a](#) arise from a 180° phase shift of the oscillation between adjacent domains. As illustrated schematically in [figure 5.5b](#), the same eigenmode of fixed handedness persists across domains; however, reversal of the Néel order imposes a phase shift in the precession of the net magnetization, resulting in opposite tr-MOKE signals. This scenario would imply a real-space manifestation of chiral antiferromagnetic oscillations with fixed handedness, where the phase of the dynamic response directly encodes the Néel vector orientation.

5.3.2 Phase-resolved tr-MOKE imaging of domain oscillations

To further investigate the phase of antiferromagnetic domain dynamics, whether excited resonantly or non-resonantly by the current-induced SOT, we extended the tr-MOKE measurements to a phase-resolved configuration. In this approach, the tr-MOKE response is probed at multiple relative delay (δ) ranging from 0° to 360° between the microwave excitation current and the probing laser pulse repetition rate. [Figure 5.6](#) and [figure 5.7](#) present the measured tr-MOKE signals at various phases across this range for two distinct cases: one where the saturating magnetic field prior to measurement was applied parallel to the current-carrying



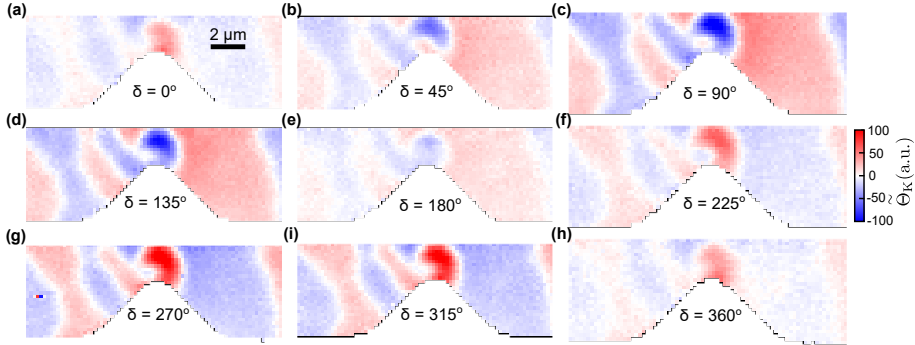


Figure 5.6: (a-h) Phase-resolved tr-MOKE images acquired at zero magnetic field under 4 GHz microwave excitation, with the relative delay (δ) between the microwave signal and laser repetition rate varied from 0° to 360° in steps. This enables pixel-by-pixel sinusoidal fitting of the Kerr signal amplitude ($\tilde{\Theta}_K$) as a function of phase delay, allowing extraction of both the absolute magnitude (Θ_K) and phase of the local domain oscillations. A saturating magnetic field was applied along the bar axis ($\alpha = 0^\circ$) prior to the measurement.

bar direction ($\alpha = 0^\circ$), and another where the field was applied perpendicular to the bar direction ($\alpha = 90^\circ$). To extract the local amplitude and phase of the oscillatory dynamics, we perform a sinusoidal fit to the tr-MOKE signal at each pixel using the measured response as a function of δ as shown in [figure 5.8](#).

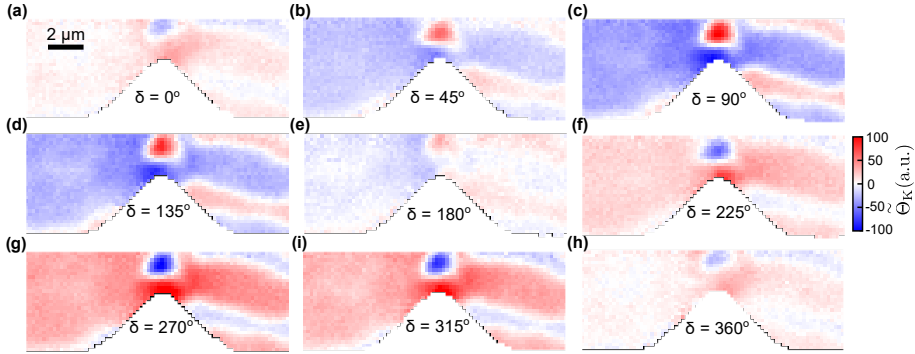


Figure 5.7: (a-h) Phase-resolved tr-MOKE images acquired at zero magnetic field under 4 GHz microwave excitation, with the relative delay (δ) between the microwave signal and laser repetition rate varied from 0° to 360° in steps. A saturating magnetic field was applied perpendicular to the bar axis ($\alpha = 90^\circ$) prior to the measurement.

The resulting amplitude and phase maps are shown in [figure 5.9](#). In [figure 5.9a](#) and [figure 5.9c](#), the amplitude maps reveal dark, line-shaped regions with near-zero signal, corresponding to DWs, while the constriction shows enhanced amplitude, consistent with locally elevated current density and stronger DL SOT-induced canting of the net magnetization. In [figure 5.9b](#) and [figure 5.9d](#), the corresponding phase maps exhibit a pronounced 180° phase shift across adjacent domains separated by the DWs; these phase-shifted regions manifest as oppositely signed red and blue contrasts in the tr-MOKE images recorded at a single delay (δ) (e.g., [figure 5.3a](#)). This robust phase signal demonstrates that the dynamic m_z component induced by the staggered DL spin-orbit field exhibits a 180° phase shift when the Néel vector reverses between adjacent antiferromagnetic domains, as described in the preceding section.

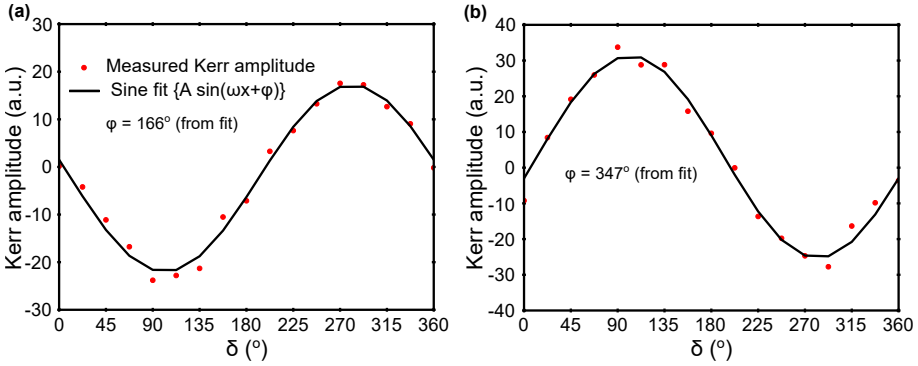


Figure 5.8: (a) and (b) tr-MOKE signal amplitude at a single pixel plotted against the relative delay (δ) between the microwave excitation and the laser repetition rate (data from [figure 5.6](#)), shown for two distinct locations within the micro-constriction device. The sinusoidal fits (solid lines) reveal that the phases of the oscillations of the net magnetization, and thus the Néel vector, differ by 180° between these two pixels, reflecting their origin from domains with reversed Néel order.

Moreover, we observed a striking directional dependence of the domain and wall orientations on the saturating magnetic field direction applied prior to the measurements: when the field is aligned parallel to the bar, domains (and walls) align perpendicular to the bar; when the field is perpendicular, domain structures realign along the bar ([figure 5.9](#)). Such reorientations can be traced to a spin-flop transition at very low fields (≈ 0 T): the Néel vector flips to lie perpendicular to the applied field to minimize the Zeeman energy, and DWs align accordingly to

reduce magnetic charges.

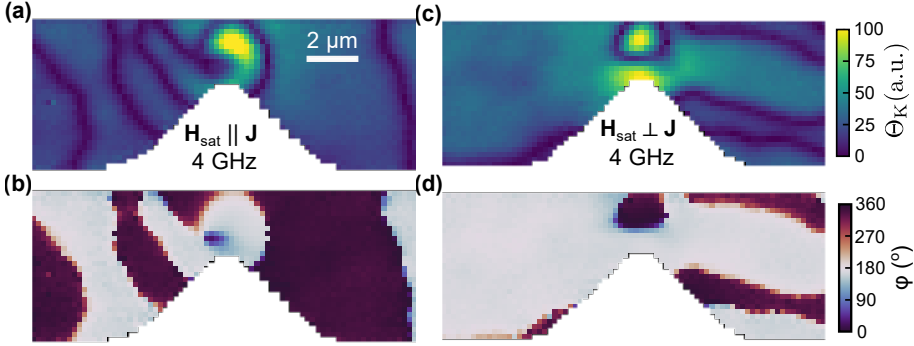


Figure 5.9: Phase-resolved tr-MOKE images at 4 GHz and zero magnetic field after saturating the sample in two orientations. (a) Magnitude map with the saturating field parallel to the bar ($\alpha = 0^\circ$). (b) Phase map corresponding to ($\alpha = 0^\circ$). (c) Magnitude map with the saturating field perpendicular to the bar ($\alpha = 90^\circ$). (d) Phase map corresponding to the ($\alpha = 90^\circ$) data. The phase- and magnitude maps are obtained from phase-resolved scans at 4 GHz in which the relative phase between the microwave drive and the laser repetition rate is swept from 0° to 360° , followed by pixel-wise sinusoidal fitting to extract the local magnetization oscillation magnitude and phase.

5.3.3 DC current driven domain-wall motion

After establishing the report of DL spin-orbit field enabled p-MOKE imaging of antiferromagnetic domains in the in-plane SAF, we now turn our attention to the investigation of SOT-driven DW motion. Previous studies have demonstrated that DWs in out-of-plane SAFs with PMA can reach exceptionally high velocities under current-driven SOT excitation. These dynamics were detected via a finite static p-MOKE signal, arising from the small uncompensated net magnetization caused by the inequality of the ferromagnetic layers forming the SAF³². However, to our knowledge, SOT-induced reversible motion of DWs in an easy-plane, in-plane SAF system, and their subsequent dynamic detection using polar tr-MOKE, have not yet been demonstrated. Figure 5.10 presents the tr-MOKE scans of the entire SAF device, acquired under continuous 4 GHz microwave excitation while systematically applying DC currents of alternating polarity. The sequence begins with a current of -14 mA, after which a tr-MOKE image is recorded in

the presence of the applied DC current. The current is then reversed to +14 mA and imaged again. This alternating procedure is repeated with progressively decreasing current magnitudes: −12 mA, +12 mA, −10 mA, +10 mA, and so forth, continuing down to −4 mA and +4 mA. At each step, the resulting domain configuration is captured via tr-MOKE imaging, enabling the visualization of current-induced DW motion throughout the device.

Application of a DC current of −14 mA, together with continuous 4 GHz microwave excitation, results in a predominantly blue tr-MOKE response across the SAF device, indicative of a favored Néel vector orientation. However, isolated red domains persist, particularly within the constriction region, where a pinned domain wall separates adjacent domains. Upon reversing the current polarity to +14 mA, the extended regions outside the constriction undergo a full reversal of the Néel order, now displaying a red tr-MOKE contrast of opposite sign. Inside the constriction, remnants of the blue domain persist, albeit with reduced amplitude, and a localized, unreversed domain remains pinned at the right edge. Reducing the current to −12 mA and then +12 mA reproduces the domain configurations observed at ± 14 mA, demonstrating reversible SOT-driven switching of the Néel order and DW motion. This reversibility persists down to ± 10 mA. At ± 8 mA, only partial domain switching is observed, with limited propagation from the constriction into the surrounding region. No discernible DW motion is detected at DC current amplitudes less than or equal to ± 6 mA, indicating a clear threshold for current-driven SOT dynamics in this device. The microconstriction acts as both a nucleation site and a pinning center for antiferromagnetic domains and DWs. Below 10 mA, the applied SOT is insufficient to overcome the pinning potential and initiate domain reversal or wall displacement.

As has been introduced earlier [section 2.3.2](#), SOT consists of DL and FL components^{124,126,214}. In SAFs and AFs, the DL SOT is most effective at driving DW motion, particularly for Néel-type walls stabilized by interfacial anisotropies or DMI²³³. The ability of the DL staggered spin-orbit field to detect SAF domains, together with the observation that domain switching occurs only above a threshold current and that DW motion is reversible, suggests that the DL torque is likely the driving mechanism in our system. This also reflects deterministic SOT-driven wall dynamics rather than stochastic thermal effects. Unlike prior studies focused on static or pulsed current-driven switching and DW motion^{18,32}, our experiment uniquely combines spin-orbit field enabled p-MOKE imaging of in-plane



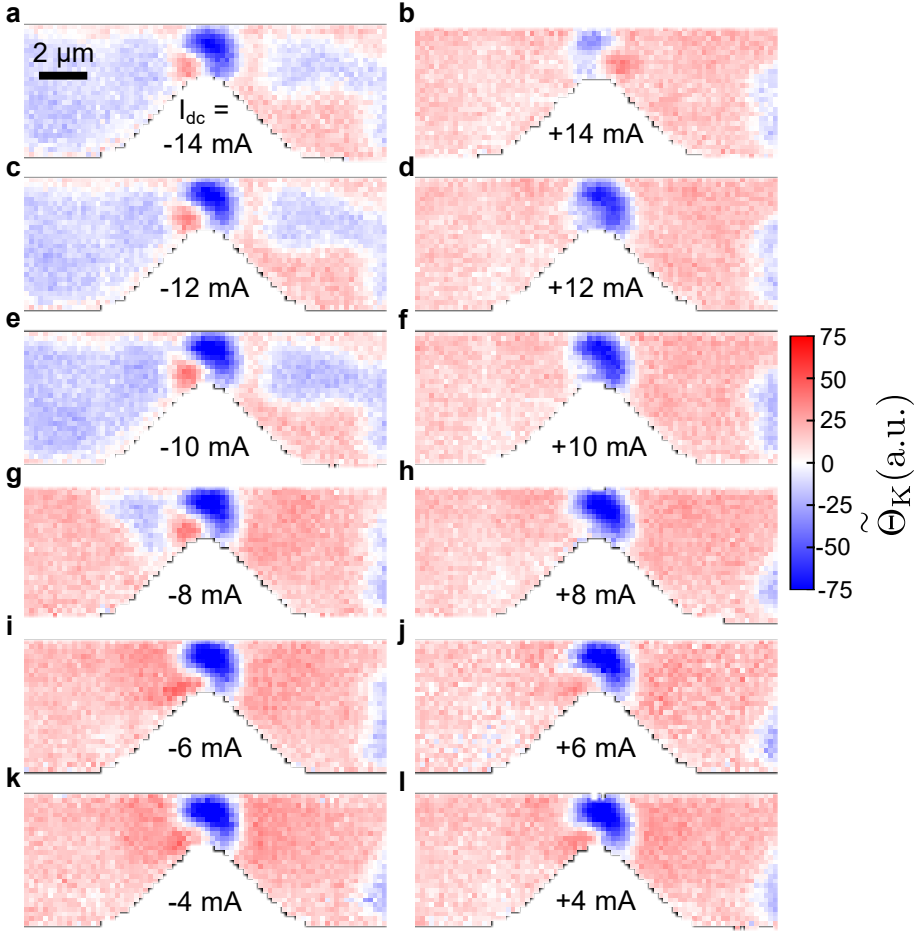


Figure 5.10: (a-l), tr-MOKE images acquired under 4 GHz microwave excitation and zero external magnetic field, while applying a DC current of decreasing amplitude with alternating polarity. The DC current sequence followed: -14 mA, +14 mA, -12 mA, +12 mA, continuing down to -4 mA and +4 mA. Prior to establishing the zero-field condition, a saturating out-of-plane magnetic field was applied perpendicular to the bar ($\alpha = 90^\circ$). The tr-MOKE signal probes the dynamic out-of-plane component (m_z) of the net magnetization. The images reveal reversible DW motion driven by SOT, dependent on the current polarity, and show a clear threshold-like behavior in the DC current amplitude required to induce DW displacement.

antiferromagnetic domains and DWs dynamics under 4 GHz drive and their displacement with gradually increasing DC bias. The micro-constriction acts as a micro-domain nucleator and pinning center, revealing how geometry crucially influences SOT effectiveness in antiferromagnetic devices.

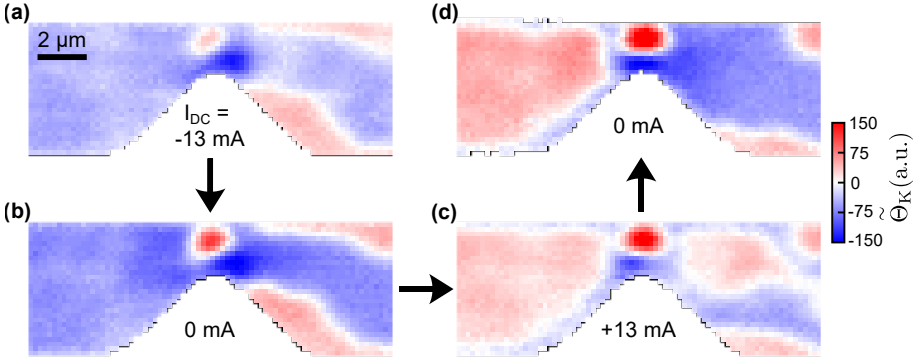


Figure 5.11: (a-d) TR-MOKE images acquired under 4 GHz microwave excitation and zero external magnetic field, while applying a DC current in the sequence of -13 mA, 0 mA, +13 mA and 0 mA respectively. These scans are performed at a fixed relative delay of 90° between the the microwave drive and the laser repetition rate.

To further highlight the deterministic and predominantly DL character of the SOT-driven dynamics, we examine the stability and reversibility of the domain configuration under intermediate current amplitudes of ± 13 mA. As shown in [figure 5.11a](#), upon applying $I_{dc} = -13$ mA together with the 4 GHz RF excitation, the device evolves into a predominantly blue domain state, while a pinned DW and mixed red-blue domains remain within the constriction. Additionally, small red domains persist outside the constricted region, reflecting local pinning and partial domain reversal. Importantly, when the DC bias is removed ($I_{dc} = 0$ mA), the domain configuration established at -13 mA is retained, demonstrating the stability of the SOT-written state in the absence of current ([figure 5.11b](#)). Reversing the polarity to $+13$ mA drives the system toward a predominantly red domain configuration across the device ([figure 5.11c](#)), again with incomplete switching and residual pinned regions. This indicates that only a current of comparable magnitude and opposite polarity can reverse the previously set magnetic state, confirming the reversible nature of the DW motion. After removing the $+13$ mA bias, a partial relaxation of some polarized red domains is observed; nevertheless,

red remains the dominant orientation ([figure 5.11d](#)). In all cases at ± 13 mA, switching remains partial rather than complete, underscoring the role of pinning and confirming that reversible DL SOT governs the controlled yet threshold-limited DW displacement in the SAF.

5.3.4 Observation of a localized mode inside the constriction

We now turn to the resonant excitation of eigenmodes in our in-plane SAF system. When an in-plane magnetic field is applied to the SAF structure, which exhibits easy-plane magnetic anisotropy, a finite net magnetization is induced along the field direction, while the Néel vector aligns perpendicular to it. Under these conditions, two linearly polarized eigenmodes can be excited. The acoustic mode corresponds to a precession of the field-induced net magnetization around the applied magnetic field at constant magnitude (see [section 3.3.2](#)). In contrast, the optical mode involves a periodic modulation of the magnetization amplitude while its orientation remains fixed along the field direction.

To probe these excitations, we performed a field-sweep measurement at a fixed microwave excitation frequency of 4 GHz. The probing laser spot was positioned at a defined location inside the micron-sized constriction of the SAF device, and the corresponding tr-MOKE response was recorded ([figure 5.12a](#)). A clear resonance behavior was observed, with resonance fields at approximately ± 26 mT. The magnetic-field dependence of the acoustic and optical eigenmode frequencies, measured electrically via spin-rectification effects in a Hall bar device (see [chapter VI](#)), is shown in [figure 5.12b](#). Notably, the observed 4 GHz resonance at ± 26 mT does not coincide with either the acoustic or optical eigenmode of the SAF system. This indicates that the detected resonance originates from a different, spatially confined excitation. The localized mode is not observed in the electrical measurements, likely due to the absence of a constriction geometry in the Hall bar device. The presence of multiple domains state at the constriction can lead to the cancellation of the electrical response. Because electrical measurements probe a global response, this cancellation can obscure the signal, whereas spatially resolved tr-MOKE can directly detect such localized modes.

To clarify its nature, we performed spatially resolved tr-MOKE imaging of the entire SAF device under continuous 4 GHz excitation at the resonant field

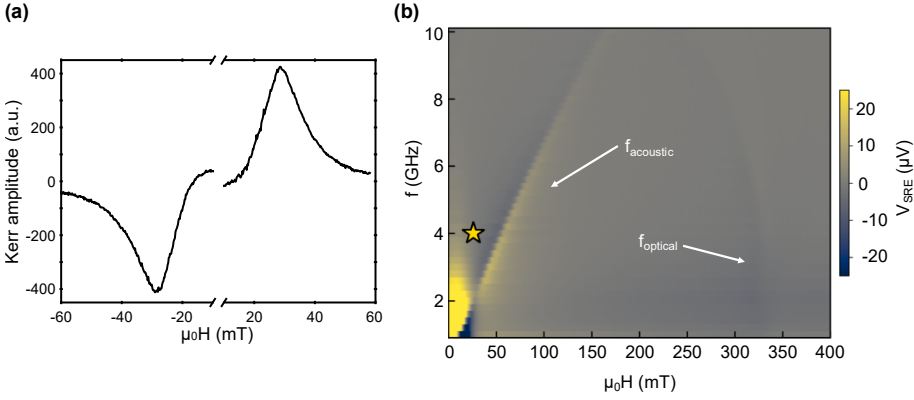


Figure 5.12: (a) Field-sweep tr-MOKE measurement performed at a fixed microwave excitation frequency of 4 GHz with the probing laser spot positioned inside the micron-sized constriction of the in-plane SAF device. A clear resonance is observed at approximately ± 26 mT. (b) Magnetic-field dependence of the acoustic and optical eigenmode frequencies of the SAF system, measured electrically via spin-rectification effects in a Hall bar device (see [chapter VI](#) for more details on SRE). The resonance observed at ± 26 mT in panel (a) (star marked symbol in panel (b)) does not coincide with either the acoustic or optical eigenmode branches and that possible electrically probed, indicating a distinct, localized excitation.

value (26 mT). The in-plane magnetic field was applied either parallel ($\alpha = 0^\circ$) or perpendicular ($\alpha = 90^\circ$) to the current-carrying bar. The resulting spatial maps are presented in [figure 5.13](#). [Figure 5.13a](#) and [figure 5.13b](#) show the tr-MOKE response at a fixed phase delay of $\delta = 0^\circ$ between the microwave excitation and the laser pulse for $\alpha = 0^\circ$ and $\alpha = 90^\circ$, respectively. The corresponding amplitude maps, obtained from pixel-by-pixel sinusoidal fits to the phase-resolved data (see [section 2.3.2](#)), are displayed in [figure 5.13c](#) and [figure 5.13d](#).

For $\alpha = 0^\circ$, a pronounced tr-MOKE signal is observed, strongly localized within the micron-sized constriction ([figure 5.13a](#)). The corresponding amplitude map confirms a substantial enhancement of the excitation amplitude inside the constriction compared to the surrounding regions. This behavior is consistent with the resonant excitation of a localized mode confined to the constricted area. In contrast, for the orthogonal field orientation ($\alpha = 90^\circ$), both the phase-resolved signal and the extracted amplitude maps exhibit only a weak response within the constriction. The strong angular dependence thus further supports the interpretation of spin-flop transition induced reorientation of the Néel order and

a geometry-assisted localized excitation mechanism, however the exact origin for this needs to be investigated.

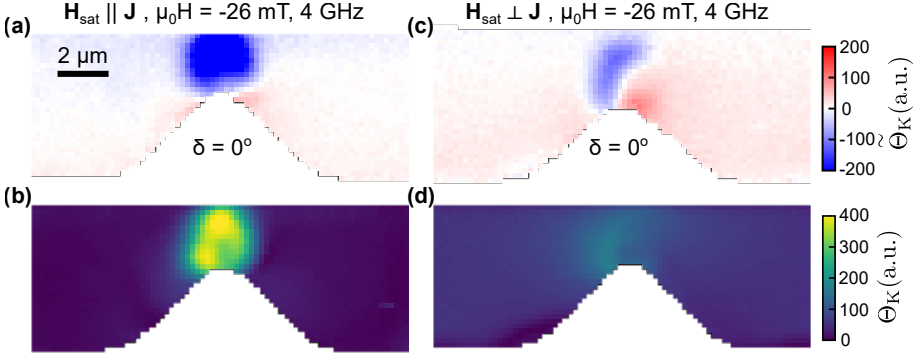


Figure 5.13: Spatially resolved tr-MOKE imaging of the SAF device under continuous 4 GHz microwave excitation at the resonant field value (26 mT). (a) tr-MOKE signal recorded at a fixed phase delay $\delta = 0^\circ$ for an in-plane magnetic field applied parallel to the current direction ($\alpha = 0^\circ$). A strong signal localized within the micron-sized constriction is observed. (b) Corresponding amplitude map extracted from phase-resolved measurements by pixel-by-pixel sinusoidal fitting (see [section 2.3.2](#)), showing a pronounced enhancement of the excitation amplitude within the constriction. (c) tr-MOKE signal for the orthogonal field orientation ($\alpha = 90^\circ$). (d) Corresponding amplitude map for $\alpha = 90^\circ$, revealing only a weak response within the constriction. The strong localization and angular dependence indicate the resonant excitation of a geometrically confined mode supported by the constriction.

The observed spatial confinement of the mode can be understood in analogy to previous studies in ferromagnetic nanostructures, where geometric confinement was shown to localize spin-wave eigenmodes. In our SAF device, the constriction modifies the local demagnetizing fields and enhances the current density, thereby increasing the local SOT and effectively reducing the damping threshold for excitation. Consequently, coherent excitation conditions are preferentially fulfilled within the constricted region, leading to spatial localization of the mode.

These results underscore the decisive role of device geometry in tailoring and controlling dynamic responses in antiferromagnetically coupled systems. Beyond acting as a nucleation and pinning site for domains (as discussed in [section 5.3.3](#)), the constriction effectively serves as a resonant cavity that supports and stabilizes localized eigenmode excitations.

5.4 Conclusion

In this chapter, we have demonstrated direct, real-space imaging of current-induced antiferromagnetic domain dynamics in an in-plane SAF using SOT-enabled polar tr-MOKE microscopy. By exploiting the staggered DL spin-orbit field acting oppositely on the two ferromagnetic layers, we overcame the intrinsic insensitivity of p-MOKE to in-plane compensated order and established a robust optical contrast mechanism for antiferromagnetic domains. Phase-resolved measurements revealed that adjacent domains with opposite Néel vector orientations oscillate with a 180° phase shift of the dynamic out-of-plane magnetization component. This phase contrast directly encodes the Néel order and confirms that the detected signal originates from staggered SOT-driven dynamics rather than from static uncompensated moments. The observed frequency dependence further suggests the presence of a weak in-plane uniaxial anisotropy, enabling resonance-like excitation behavior and potentially chiral eigenmode dynamics in the easy-plane SAF.

Beyond domain imaging, we demonstrated reversible, threshold-like DC current-driven DW motion under continuous microwave excitation. The deterministic polarity dependence and stability of the written states identify the DL SOT as the possible driving mechanism. The microconstriction plays a dual role as a nucleation and pinning center for domains and as a region of enhanced current density, thereby strongly influencing the efficiency of SOT-induced switching.

Finally, we identified a geometry-assisted, spatially localized excitation mode confined to the constriction region under applied magnetic field. Its strong angular dependence and deviation from the acoustic and optical eigenmodes highlight the decisive role of device geometry in shaping antiferromagnetic spin dynamics.

Overall, this work establishes SOT-enabled tr-MOKE microscopy as a powerful tool for probing and controlling dynamic Néel order in compensated in-plane SAFs. The ability to optically resolve domain oscillations, reversible DW motion, and localized eigenmodes in real space provides an essential experimental platform for the development of high-speed, geometry-engineered antiferromagnetic spintronic devices.



VI

SOT-driven chiral self-oscillations in SAF

In this chapter, we present a SOT-driven oscillator realized in a nanoconstriction patterned from an easy-plane SAF. Spin rectification measurements reveal current-driven excitations of both optical (OM) and acoustic (AM) modes, along with additional resonances emerging beyond a threshold DC current near the spin-flop transition. These anomalous resonances are consistent with self-sustained, chiral oscillations whose handedness is governed by the DC current polarity and which can injection-lock to an external microwave frequency. Complementary macrospin and micromagnetic simulations support the existence of these chiral self-oscillations and further predict complex non-linear dynamics including chaotic behavior near the spin-flop transition. Experimentally, the onset of multiple frequency-insensitive resonance modes above the threshold current provides qualitative signatures of these complex dynamics.

Keywords: spin-orbit torques (SOTs), antiferromagnetic nano-oscillators, spin rectification effect, chiral self-oscillations, coherent antiferromagnetic spintronics.

This chapter is based on: **Rout, P.K.**, Godinho, J. et al. Spin-orbit torque driven antiferromagnetic oscillator. Under review.

6.1 Introduction

STT and SOT nano-oscillators represent a rapidly advancing frontier in spintronics and magnonics, offering new possibilities for compact, energy-efficient, and high-speed signal generation^{73,88}. Their potential applications extend across neuromorphic computing, microwave-assisted magnetic recording, magnonics, and next-generation wireless communication technologies^{56,57,73,88,91-94,198,199}. Owing to their highly tunable operation, compatibility with standard CMOS platforms, and ability to mutually synchronize, these nanoscale oscillators have become a cornerstone in the pursuit of functional spintronic devices^{57,92-94,234}.

STT nano-oscillators are typically implemented using spin-valve or MTJ geometries, where a spin-polarized current transfers angular momentum to the free magnetic layer^{73-78,119}. This process compensates intrinsic damping, thereby sustaining steady-state precession of the magnetization at microwave frequencies. In contrast, SOT nano-oscillators utilize charge-to-spin conversion enabled by strong SOC in adjacent HM layers to excite similar magnetization dynamics, offering an alternative pathway to current-induced oscillations without relying on spin filtering through a magnetic polarizer⁷⁹⁻⁸⁴. In both cases, the oscillation frequency is predominantly determined by the strength of the applied magnetic field, while the precession sense, or intrinsic chirality, is dictated by the orientation of the $\mathbf{M} \times \mathbf{H}$ torque, yielding a fixed right-handed rotation. Therefore, in ferromagnetic systems, both STT and SOT oscillators exhibit inherent constraints: for a given magnetic field direction, the chirality of precession cannot be reversed, and frequency tunability remains limited, typically achievable only through modest adjustments such as voltage-controlled magnetic anisotropy^{235,236}.

This limitation motivates the exploration of antiferromagnetic and synthetic antiferromagnetic systems, where the absence of net magnetization and stray fields and the presence of exchange-enhanced spin dynamics open entirely new regimes of oscillatory behavior^{95,96,187,188}. These systems provide highly promising avenues for spintronic oscillators, combining ultrafast dynamics that can push operational frequencies into the terahertz range with bias-current control over both frequency and chirality. Theoretical studies have revealed a rich spectrum of dynamical behaviors in such systems, including chiral self-oscillations, resonantly driven left- and right-handed circularly polarized modes in easy-axis

antiferromagnets, and linearly polarized acoustic and optical modes in easy-plane systems^{188,237–240} (for discussion about eigenmodes in AFs see [section 3.3](#)). Furthermore, the frequency range of current-induced self-oscillations is predicted to be significantly broader than in conventional ferromagnetic oscillators. For instance, an insulating AF such as NiO, driven by spin currents from adjacent nonmagnetic metals, is predicted to exhibit chiral, self-sustained Néel-order oscillations spanning the frequency gap between the acoustic and optical eigenmodes²³⁷. Important to note here that the direction of the driving current directly controls the chirality of such antiferromagnetic excitation. Another promising platform is metallic $\mathcal{PT}$ -symmetric AFs, where SOT can directly induce and control self-oscillations in a single AF layer exhibiting finite dissipative SOT²³⁸. In these conductive systems, the dynamics and handedness of oscillations can be probed experimentally through phase-sensitive nonlinear rectification, arising from time-dependent magnetotransport effects such as AMR or the AHE, particularly in noncollinear antiferromagnets^{241,242}.

Beyond periodic oscillations, antiferromagnetic systems with easy-axis anisotropy are predicted to display chaotic dynamics when driven near the spin-flop transition^{239,240}. In this case the geometry of the studied system is represented by a pillar structure that consists of three layers: ferromagnetic polarizer, nonmagnetic spacer, and the studied AF free layer with easy-axis anisotropy out of the plane ([figure 6.1](#)). Two stable static ground state solutions, below (A) and above (B) the spin-flop transition, are shown in [figure 6.1a](#) and [figure 6.1b](#), respectively. Their corresponding eigenfrequency spectra are plotted in [figure 6.1d](#) (dotted lines) for zero current density ($j = 0$). As expected (see [chapter III](#)), below the spin-flop transition, two chiral eigenmodes emerge: a high-frequency mode (ω_H^A) and a low-frequency mode (ω_H^B). Both modes become degenerate at zero magnetic field, while ω_H^B softens beyond the spin-flop transition. The application of STT significantly alters the eigenmode characteristics of both states (A) and (B). It lifts the degeneracy between the high- and low-frequency modes at zero field, introducing a splitting proportional to j^2 , and reduces the critical field for the spin-flop transition, where state (A) becomes unstable. Near this instability, the high- and low-frequency branches exhibit asymptotic behavior, deviating from their linear field dependence observed at $j = 0$. Additionally, even a small spin torque opens a gap $b_{c1} < b < b_{c2}$ around the spin-flop field where no stable static solutions exist. As shown in [figure 6.1e](#), this gap region is dominated by chaotic self-oscillations. This behavior, unprecedented in ferromagnetic oscillators but

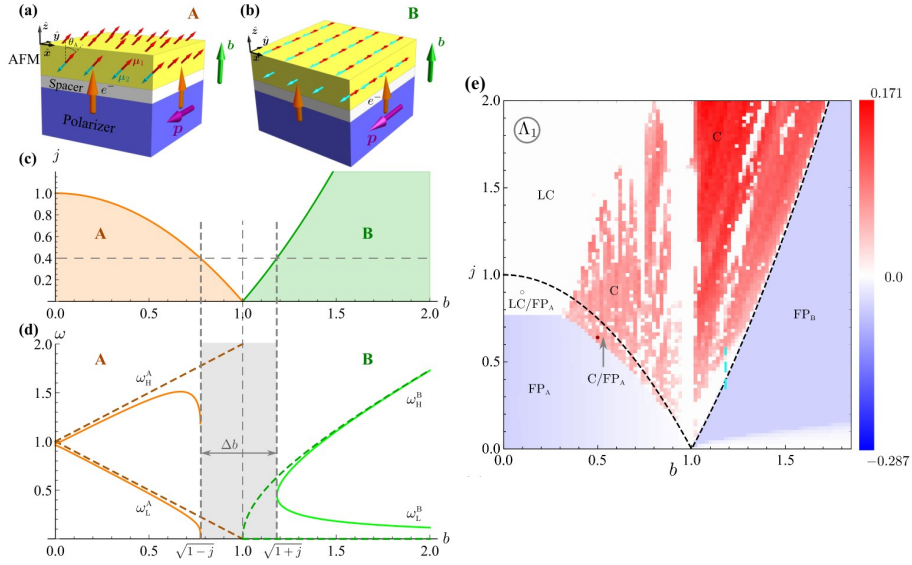


Figure 6.1: Geometry of a STT driven easy-axis AF oscillator. Stable static ground states below (A) and above (B) the spin-flop transition are shown in (a) and (b), respectively. (c) Stability domains of solutions 'A' and 'B'. The unit vector $\mathbf{p}$ indicates the spin polarization direction, $b = B/B_{\text{sf}}$ is the dimensionless magnetic field normalized to the spin-flop field, and j is the dimensionless current density. (d) Eigenfrequencies of solutions A and B for $j = 0$ (dashed) and $j = 0.4$ (solid); the shaded region of width $\Delta b \approx j$ marks a gap where no static solutions exist. (e) Lyapunov exponent map (λ_1) in parameter space for Gilbert damping $\alpha = 0.1$, showing different oscillatory dynamics: Regions in blue with $\lambda_1 < 0$ for fixed points A and B, LC ($\lambda_1 = 0$) for stable limit cycles, and C ($\lambda_1 > 0$) for chaotic dynamics. In some regions, multiple oscillation modes coexist. Dashed lines indicate the stability boundaries of static solutions, corresponding to panel (d). Reproduced from [239].

reminiscent of vortex-based systems with feedback-driven core reversal, opens exciting possibilities for applications in neuromorphic computing and reservoir computing architectures^{243,244}. However, to the best of our knowledge, despite the rich and complex dynamics predicted for AFs, no experimental study has yet reported signatures consistent with current-driven self-oscillatory behavior in antiferromagnetic systems.

In this work, we investigate both resonant eigenmodes and self-oscillatory dynamics in a SAF multilayer, which provides a structurally simple yet highly versatile platform for studying antiferromagnetic dynamics. Despite its simplicity,

the SAF preserves the essential physical ingredients necessary to emulate complex current-driven antiferromagnetic behavior (see [chapter I chapter II](#) for detailed motivation). The SAF structure consists of two ferromagnetic layers separated by a nonmagnetic spacer that mediates a RKKY interaction, establishing antiferromagnetic coupling between the layers⁴¹. The multilayer is further sandwiched between tantalum (Ta) and platinum (Pt) layers, which exhibit opposite SHAs. When a charge current flows through this stack, spin currents generated at the HM interfaces are injected into the SAF from both sides with identical spin polarization^{20,124,125,134,135,137} (see [section 2.3.2](#)). The resulting SOT acts on the SAF's antiferromagnetic order in a manner analogous to the STT predicted in collinear antiferromagnets, where chaotic self-oscillatory behavior is expected near the spin-flop transition²³⁹.

By driving the SAF with current-induced SOT, we observe both the characteristic antiferromagnetic eigenmodes as well as indications of self-sustained oscillations. Using a combination of rf current excitation and an in-plane magnetic field of variable magnitude and orientation, we detect linearly polarized acoustic and optical resonance modes via the spin rectification effect (SRE)^{219,220,241}. Additional rectification peaks appear once the applied DC current surpasses a critical threshold, signaling the onset of self-sustained chiral oscillations of the antiferromagnetic order. These oscillations become injection-locked to the frequency of the RF signal, while the polarity-dependent sign reversal of the SRE confirms that the chirality of oscillation is determined by the direction of the DC current. To substantiate these experimental findings, we performed macrospin and micromagnetic simulations incorporating realistic material parameters and device geometry. Near the spin-flop transition, the simulations reproduce current-induced self-oscillations spanning a wide frequency range between the AM and OM and demonstrate injection locking to the external rf drive, consistent with experimental observations. Moreover, Lyapunov exponent analysis and bifurcation mapping of the simulated dynamics reveal the onset of chaotic behavior in the DC-driven self-oscillations of the SAF, corroborating the predicted nonlinear character of SOT-driven antiferromagnetic dynamics. Important to note here, throughout this work, DC-current driven auto-oscillatory dynamics are inferred from anomalous spin-rectification signatures and injection-locking behavior, and are supported by micromagnetic and macrospin simulations, rather than by direct detection of microwave emission.

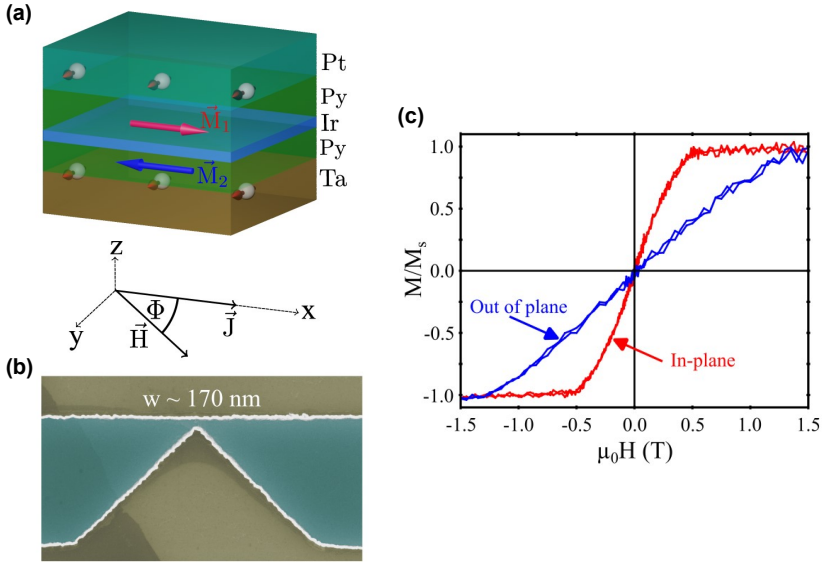


Figure 6.2: (a) Schematic of the SAF multilayer comprising two permalloy (Py: $\text{Ni}_{0.81}\text{Fe}_{0.19}$) layers antiferromagnetically coupled via a 0.34 nm Ir spacer, sandwiched between Pt and Ta layers. The green arrows indicate the spin-polarization directions injected from the heavy-metal interfaces when a charge current flows through the multilayer stack. Red and blue arrows represent the magnetization vectors ($\mathbf{M}_1$) and ($\mathbf{M}_2$) of the upper and lower Py layers, respectively. The in-plane magnetic field ($\mathbf{H}$) and charge current ($\mathbf{J}$) are applied at an angle Φ . (b) Scanning electron micrograph of a fabricated nanoconstriction device with a constriction width (w) of approximately 170 nm. (c) Normalized magnetization as a function of the magnetic field H for in-plane (red) and out of plane (blue) directions measured in an unpatterned SAF.

6.2 Experimental Methods

6.2.1 Sample and device fabrication

The SOT-driven antiferromagnetic oscillator is based on a nominally compensated in-plane SAF with the following structure: Si/SiO₂ (substrate)/Ta (5 nm)/Py (2 nm)/Ir (0.34 nm)/Py (2 nm)/Pt (5 nm). The SAF multilayer stack (figure 6.2a) consists of two 2 nm-thick permalloy (Py: $\text{Ni}_{0.81}\text{Fe}_{0.19}$) layers that are antiferromagnetically coupled through a 0.34 nm Ir spacer via the RKKY interaction. The

detailed sample preparation and nano-constriction device fabrication procedures are identical to those described in [section 5.2.1](#). To enhance the current density and spatially confine the magnetic excitations, the multilayer stack was patterned into a bar featuring a single notch-shaped nanoconstriction ([figure 6.2b](#)). Vibrating sample magnetometry (VSM) measurements reveal a predominant easy-plane magnetic anisotropy ([figure 6.3c](#)). Additionally, a weak in-plane biaxial anisotropy is induced by the nanofabrication process which will be discussed later.

6.2.2 Spin rectification effect

In [chapter III](#), we discussed field and current-induced magnetization dynamics in FMs and AFs in the context of FMR and AFMR. According to the LLG equation, the magnetization $\mathbf{M}$ or Néel vector $\mathbf{L}$ precesses around the effective field $\mathbf{H}_{\text{eff}}$, with the precession amplitude decaying over time due to Gilbert damping. External energy supplied through an alternating magnetic field $\mathbf{h}_{\text{RF}}$ can drive precession, leading to resonance when the driving frequency matches the natural precession frequency. Electrical excitation of FMR became possible following the discovery of STT^{39,119,120}, where an rf spin-polarized current drives magnetization dynamics via DL (τ_{DL}) and FL (τ_{FL}) torques. In HM/FM bilayers, charge-to-spin conversion through the SHE or iSGE enables SOT-driven FMR ([section 2.3.2](#)). SOT-driven FMR has since been demonstrated in Pt/NiFe and (Ga,Mn)As systems, establishing an efficient, all-electrical pathway for exciting magnetization dynamics without the need for a separate polarizer^{219,220}. In these systems, an rf charge current generates a spin current that exerts DL (τ_{DL}) and FL (τ_{FL}) torques, alongside an Oersted field torque (τ_{Oe}), as described in [equation \(3.33\)](#). As illustrated in [figure 6.3](#) for Pt/NiFe bilayers, charge-to-spin conversion within the HM layer produces a non-equilibrium spin accumulation at the interface. This accumulation diffuses into the FM, exerting SOTs on the local magnetization. Because these driving torques oscillate at the rf frequency, they sustain a coherent precession of the magnetization about the effective field, resulting in an oscillatory AMR.

The dynamics are detected electrically via the SRE, which arises from the nonlinear mixing of an oscillating resistance (AMR in case of [figure 6.3](#)) and the rf current ($I_{\text{rf}}(t)$)²⁴¹. Magnetization precession modulates the device resistance through magnetoresistive effects ([section 2.4](#)), yielding a time-dependent

resistance $R[\mathbf{H}(t)]$. The rectified dc voltage along the current path is given by²⁴¹

$$V_{\text{SRE}} = \langle \text{Re}\{I_{\text{rf}}(t)\} \text{Re}\{R[\mathbf{H}(t)]\} \rangle = \frac{1}{2} I_0 h_0 \cdot \nabla R(\mathbf{H}_0) \cos \Phi, \quad (6.1)$$

where Φ denotes the phase difference between current and magnetization oscillations. The voltage originates from the first-order Taylor expansion of $R[\mathbf{H}(t)]$ about the static field $\mathbf{H}_0$ and averaging over one microwave cycle. In STT- and SOT-driven FMR, the resulting lineshape consists of symmetric and antisymmetric Lorentzian components. These correspond to the DL and FL torques, respectively, enabling quantitative extraction of their efficiencies^{219,220}. The SRE therefore provides a powerful all-electrical tool for probing magnetization dynamics, spin-wave excitations, nonlinear phenomena, DW motion, spin currents, and microwave imaging²⁴¹. From this description, three essential ingredients are required for spin rectification: (1) a magnetoresistive mechanism, where the device resistance depends on the magnetization orientation; (2) a torque that drives rf magnetization, generating a time-varying resistance; and (3) an rf current that couples to this dynamic resistance to produce a measurable dc voltage.

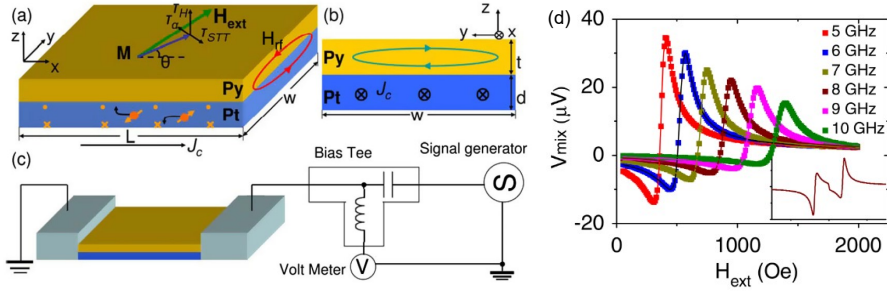


Figure 6.3: (a) Schematic of a Pt/Py bilayer thin film illustrating the STT (τ_{STT} or τ_{DL}), the Oersted-field-induced torque τ_{H} , and the damping torque τ_{α} direction. ϕ denotes the angle between the magnetization $\mathbf{M}$ and the microstrip, while $\mathbf{H}_{\text{ext}}$ is the applied external magnetic field. The SHE in Pt deflects spins (polarized in the sample plane) toward the top surface, generating a spin current incident on the Py layer. Here, τ_{STT} is analogous to the damping-like torque τ_{DL} , while the field-like torque τ_{FL} is absorbed into τ_{H} as both share the same symmetry. (b) Side view of the Pt/Py system, where the solid line indicates the Oersted field generated by current flowing solely in the Py layer, which produces no net effect on the Py AMR. (c) Schematic circuit of the ST-FMR measurement. (d) ST-FMR spectra of a Pt(6 nm)/Py(4 nm) bilayer measured at frequencies between 5 and 10 GHz. Inset: ST-FMR spectrum at 8 GHz for both positive and negative H_{ext} .

To probe the spin dynamics of the SAF nanoconstriction, we applied a rf current and measured the resulting spin rectification voltage V_{SRE} while sweeping an in-plane magnetic field. This voltage arises from nonlinear mixing between the time-dependent resistance, modulated by AMR and GMR, and the applied RF current. All-electrical SRE measurements were performed at room temperature. The SAF nanoconstriction chip was mounted on a custom rf PCB with an integrated coplanar waveguide (CPW) for high-frequency access and wire-bonded to the CPW. A radio-frequency current (I_{rf}) and a direct current (I_{DC}) were applied through the high- and low-frequency ports of a bias tee, using a signal generator and a precision source-measure unit, respectively. The assembled device was placed in a room-temperature electromagnet, with the applied field monitored via a Hall sensor and a gaussmeter. SRE measurements followed a field-sweep protocol at fixed rf frequency to identify resonance conditions. The rf current generates SOTs within the SAF nanoconstriction and, in principle, an Oersted-field torque; however, the latter is negligible due to its small magnitude and symmetry as will be explained further. Because V_{SRE} is very small, amplitude modulation of the RF current at 333 Hz and lock-in detection were employed. The rf PCB assembly was mounted on a non-magnetic piezo rotation stage to enable measurements at different in-plane field orientations.

6.2.3 Micromagnetic simulations

In order to verify the existence of the SOT-driven self-oscillations in the SAF nano-constriction, micromagnetic simulations were performed using MuMax3²¹³. For the simulations, a reduced device geometry of $1\ \mu\text{m} \times 0.5\ \mu\text{m}$ with a 70 nm constriction was used, scaled from the experimental $15\ \mu\text{m} \times 5\ \mu\text{m}$ structure while preserving the aspect ratios, including the asymmetric constriction with its 90° opening angle. Only the two Py layers of the SAF were modeled. Current and Oersted-field distributions were obtained from COMSOL simulations of the same geometry and imported into MuMax3 (figure 6.4).

The SAF stack was discretized on a $512 \times 256 \times 3$ mesh, corresponding to a cell size of approximately $2 \times 2 \times 1\ \text{nm}^3$. To suppress artefacts from magnons reflected at the wide leads, Gilbert damping was gradually increased toward the device boundaries, and periodic boundary conditions were employed to approximate the extended dimensions of the experimental device. Material parameters were set to

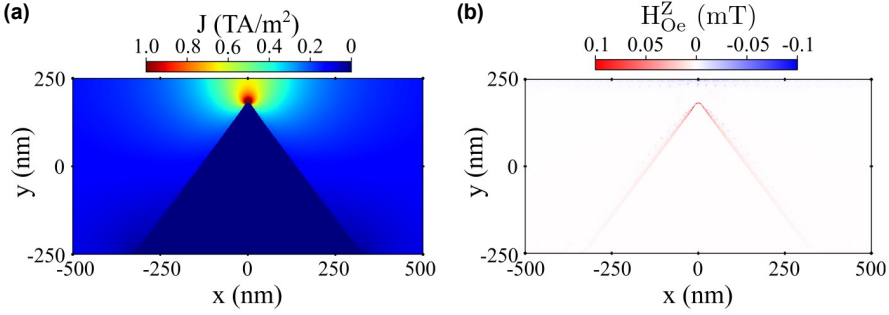


Figure 6.4: Simulated current densities and Oersted field: (a) The magnitude of the lateral current densities in Pt, when a 1 mA direct current is applied across the entire SAF stack, as simulated using COMSOL. (b), The corresponding out of plane component of the Oersted field at the Py site.

a saturation magnetization $M_s = 7.9 \times 10^5$ A/m, intra-layer exchange stiffness $A = 1.1 \times 10^{-11}$ J/m, and Gilbert damping $\alpha = 0.01$. The interlayer antiferromagnetic coupling was adjusted to reproduce the experimentally observed spin-flip fields in both in-plane and out-of-plane configurations, yielding a exchange field of 0.25 T. A spin Hall angle of 0.15 was used to model the SOT. The system was first relaxed to its ground state and then evolved under current and/or field excitation until reaching a steady-state dynamical regime.

6.3 Results and discussion

6.3.1 Spin-orbit torque excitation of antiferromagnetic eigenmodes in SAF

The SAF multilayer stack consists of two 2-nm-thick permalloy (Py: $\text{Ni}_{0.81}\text{Fe}_{0.19}$) layers antiferromagnetically coupled through a 0.34-nm Ir spacer via the RKKY interaction. The SAF is further sandwiched between a Ta seed layer and a Pt capping layer, which have opposite spin Hall angles (figure 6.2a). Consequently, when a charge current flows through the multilayer, both HM interfaces generate spin currents of identical spin polarization that are injected into the SAF from opposite sides, giving rise to SOTs that drive magnetization dynamics. We first examine

the excitation of the SAF eigenmodes, which, owing to the easy-plane anisotropy, comprise the acoustic and optical modes (see [section 3.3.2](#) and [figure 6.5a](#) for schematics). These antiferromagnetic eigenmodes are driven by SOTs from an rf current and detected through the SRE.

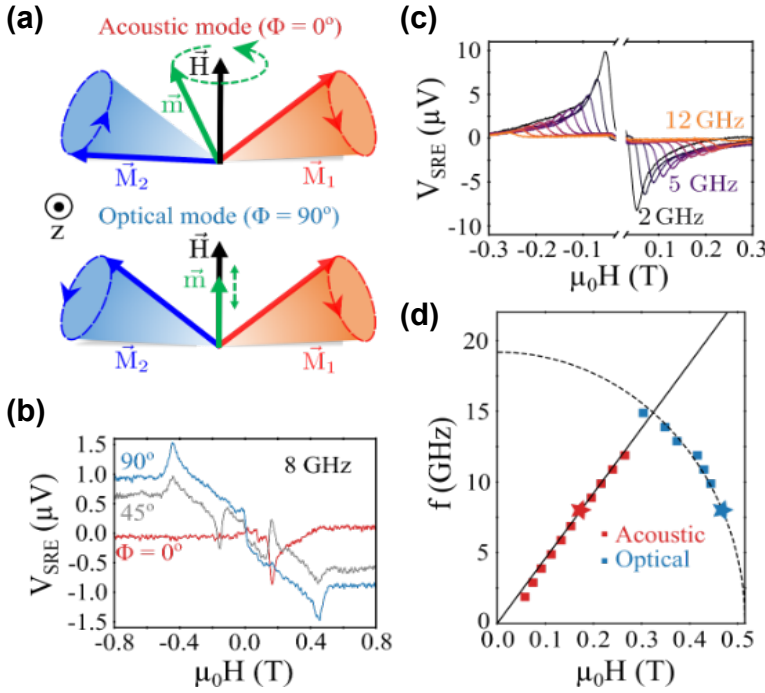


Figure 6.5: (a) Schematic of the magnetization dynamics for the acoustic mode (top, excited at $\Phi = 0^\circ$) and the optical mode (bottom, excited at $\Phi = 90^\circ$). The acoustic mode corresponds to out-of-plane precession of the net magnetization $\vec{M} = \vec{M}_1 + \vec{M}_2$, whereas the optical mode exhibits in-plane precession. The in-plane magnetic field $\vec{H}$ and charge current $\vec{J}$ are applied at an angle Φ . (b) Spin rectification voltage V_{SRE} as a function of magnetic field for $\Phi = 0^\circ$, 45° , and 90° at 8 GHz and 15 dBm rf power. Depending on Φ , either selective or simultaneous excitation of the acoustic and optical modes is observed. (c) V_{SRE} versus magnetic field for rf frequencies from 2 GHz to 12 GHz, illustrating the frequency dependence of the AM. (d) Extracted resonance frequencies of the acoustic (red symbols) and optical (blue symbols) modes as a function of magnetic field. Solid and dashed black curves show the analytical mode frequencies for easy-plane AFs, using $\mu_0 H_{\text{ex}} = 0.25$ T and $\mu_0 H_{\text{an}} = 0.98$ T, independently obtained from VSM measurements. The star marks the data point corresponding to the 8 GHz measurements in panel (b).

An in-plane magnetic field produces a net magnetization ($\mathbf{M}_1 + \mathbf{M}_2$) aligned with the field and a Néel vector ($\mathbf{M}_1 - \mathbf{M}_2$) oriented perpendicular to it. The mode excited depends on the azimuthal angle Φ between the magnetic field and the rf current, defined by the bar orientation (see [figure 6.2b](#)). In the AM, the field-induced net magnetization precesses at constant magnitude around the effective field direction ([figure 6.5a](#)). In contrast, in the OM, the magnitude of the net magnetization oscillates while its orientation remains fixed along the field ([figure 6.5a](#)). As illustrated in [figure 6.5a](#), the acoustic and optical modes can be selectively or jointly excited depending on Φ , consistent with the symmetry of the SOT excitation. Because the SOT field is oriented perpendicular to the rf current, exclusive excitation of one mode is possible by applying the magnetic field either collinear or perpendicular to the current direction. Specifically, the AM is driven when the field is aligned with the bar ($\Phi = 0^\circ$), while the OM is selectively excited when the field is perpendicular ($\Phi = 90^\circ$). At an intermediate orientation ($\Phi = 45^\circ$), both modes can be simultaneously excited.

[Figure 6.5b](#) shows distinct resonance peaks at 8 GHz corresponding to the optical and acoustic eigenmodes, originating from in-plane and out-of-plane oscillations of the net magnetization, respectively. When the field is applied parallel to the current ($\Phi = 0^\circ$), only the AM is observed, appearing as a resonance at positive field polarity. The absence of a corresponding resonance at negative polarity is attributed to a pinned 180° domain wall at the nanoconstriction (as shown from the tr-MOKE measurements in [chapter V](#)), which suppresses SRE contributions from oppositely magnetized domains. A slight misalignment (e.g., $\Phi = 5^\circ$) restores acoustic resonances at both polarities ([figure 6.5c](#)), indicating a uniform domain configuration. At $\Phi = 45^\circ$, both modes emerge, and at $\Phi = 90^\circ$ only the OM is present.

In our SAF device, the OM can be excited only by a non-staggered DL torque. Its SRE voltage arises predominantly from a GMR-type contribution, since the two FM sublayers precess out of phase and periodically modulate their relative alignment. By contrast, AMR contributions are negligible for this mode. Our dc magnetotransport measurements, as will be discussed later in this chapter, show that the magnitudes of the GMR- and AMR-related resistances in the device are comparable, ensuring sufficient sensitivity to this mode. By symmetry, the GMR-based SRE voltage generated by a DL torque must reverse sign when the applied magnetic field is reversed. This behavior is exactly what we observe for the

optical mode (figure 6.5b). Since FL torques or Oersted fields cannot excite this mode in our SAF geometry, the observed sign inversion provides unambiguous evidence that the optical mode is driven by a DL-SOT.

The AM in our SAF system can, in principle, be excited by both DL and FL (including Oersted) torques. Its SRE signal is dominated by an AMR-type contribution, because in this mode the two sublayer magnetizations precess largely in phase, keeping their relative alignment nearly constant such that $\frac{dR_{\text{GMR}}}{dt} \approx 0$. If the AM were excited predominantly by a FL torque or by the Oersted field, the resulting AMR-type rectified voltage would be symmetric with respect to magnetic-field polarity and would not change sign upon field reversal. In contrast, for excitation by a DL torque, the sign of the AM SRE voltage depends on whether the local SAF domain at the nanoconstriction reverses when the external field is inverted. If the antiferromagnetic order flips, the SRE voltage reverses sign; if it does not, the voltage retains the same sign. This behavior is fully analogous to the well-established response of ferromagnetic Kittel modes driven by DL spin Hall torques in conventional SOT-FMR experiments, where symmetric voltage components reverse sign with magnetization reversal, whereas pure Oersted-field excitation does not²²⁰. In our measurements, we indeed observe both situations. For example, in figure 6.5b, the AM SRE signal is suppressed near negative fields, consistent with the presence of two oppositely oriented SAF domains separated by a 180° DW pinned at the constriction, whose contributions from the two domains cancel. This confirms that the polarity of the AM SRE signal is governed by the local domain configuration and is consistent with excitation by a DL torque.

Furthermore, the opposite polarities of the acoustic and optical resonance peaks at a given field sign indicate distinct rectification channels, AMR for the acoustic mode and GMR for the optical mode. As will be shown below in this chapter, we experimentally determine that in our device the AMR and GMR contributions are comparable in magnitude but opposite in sign. Since the acoustic and optical modes are dominated by AMR and GMR, respectively, this naturally results in opposite signs of V_{SRE} for the two modes.

At higher magnetic fields, Joule heating can introduce additional contributions to the detected signal through magneto-thermoelectric effects. In particular, a background voltage offset arises from the ANE, generated by a vertical temperature gradient together with a net magnetization component perpendicular to the bar. Such a perpendicular magnetization component develops when the applied

magnetic field tilts the two antiferromagnetically coupled layers, as observed for $\Phi = 90^\circ$ and $\Phi = 45^\circ$ in [figure 6.5b](#).

As has been discussed before, SRE line shapes result from mixing between the rf current and a time-dependent resistance, and they can be expressed as a superposition of symmetric and antisymmetric Lorentzian contributions, with weights determined by the relative roles of DL torque, FL torque/Oersted fields, and phase shifts in the measurement chain. In the [chapter A](#), a concise analytical discussion has been given for our SAF geometry, explicitly deriving the dominant rectification channels, AMR-type rectification for the AM and GMR-type rectification for the OM. However, in our measurements, accurately separating the background response from the Lorentzian resonance signal remains challenging. This difficulty arises from the simultaneous variations in the AMR, GMR, and ANE responses during the magnetic field sweep, as well as from additional contributions originating from multiple magnetic domains at the constriction.

[Figure 6.5c](#) presents the frequency dependence of the AM, showing that the resonance shifts to higher magnetic fields with increasing rf frequency, consistent with theoretical expectations. In [figure 6.5d](#), the experimentally extracted resonance fields are compared with the analytical mode frequencies for easy-plane AFs (see [chapter III](#)). The acoustic mode follows $f_{\text{acoustic}} = \gamma H \left(1 + \frac{|H_A|}{2H_E}\right)^{1/2}$ while the optical mode obeys $f_{\text{optical}} = \gamma \left(2H_E|H_A| - \frac{|H_A|}{2H_E}H^2\right)^{1/2}$, where H_E and H_A denote the antiferromagnetic RKKY exchange field and the anisotropy field, respectively. Using independently determined values $\mu_0 H_E = 0.25$ T and $\mu_0 H_A = 0.98$ T from VSM measurements ([section 6.2.1](#)), we find excellent agreement between experiment and analytical theory.

6.3.2 Micromagnetic simulations of SOT-driven self-oscillations in SAF

To gain further insight into the current-driven spin dynamics in the SAF nanoconstriction, we performed micromagnetic simulations using the GPU-accelerated open-source solver MuMax3 (see [section 6.2.3](#))²¹³. The simulations incorporate the experimentally determined material parameters and device geometry of the SAF nanoconstriction shown in [figure 6.2](#). They reveal the onset of self-sustained

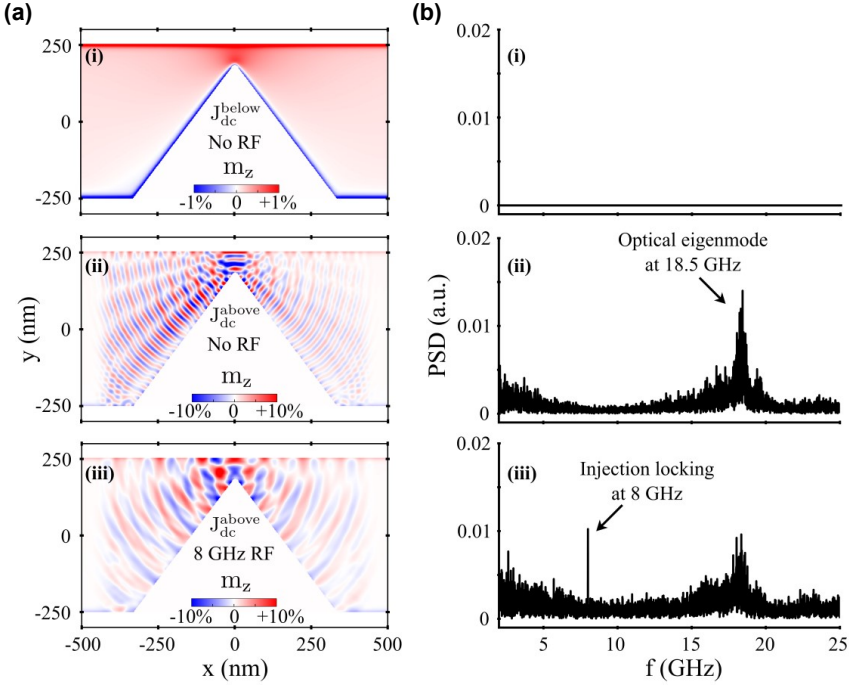


Figure 6.6: (a) Time snapshots of the normalized net magnetization for (i) J_{dc}^{below} , (ii) J_{dc}^{above} , and (iii) $J_{dc}^{above} + J_{rf}$. (b) Corresponding power spectral densities (PSD) for the three bias conditions. Here, $J_{dc}^{below} = 1.7 \times 10^8$ A/cm², $J_{dc}^{above} = 2.2 \times 10^8$ A/cm², and $J_{rf} = A \sin(2\pi f t)$ with $f = 8$ GHz and $A = 0.5 \times 10^8$ A/cm². The normalized out-of-plane component is defined as $m_z = \hat{z} \cdot (\mathbf{M}_1 + \mathbf{M}_2) / (|\mathbf{M}_1| + |\mathbf{M}_2|)$. A small in-plane magnetic field of 45 mT is applied along the bar direction during the simulations.

oscillations once the applied dc current exceeds a critical threshold. To characterize the antiferromagnetic dynamics, we analyze the time evolution of the net magnetization $\mathbf{M} = \mathbf{M}_1 + \mathbf{M}_2$, which becomes non-zero only when the two layers are dynamically canted, either during AF excitations or during the propagation of antiferromagnetic magnons.

Figure 6.6 summarizes representative results for different driving regimes at an applied in-plane magnetic field of 45 mT along the bar. For a subthreshold dc current density of $J_{dc}^{below} = 1.7 \times 10^8$ A/cm² at the constriction, the DL-SOT induces only a slight out-of-plane tilt of the magnetizations, producing a weak static net magnetization (panel (i) of figure 6.6a). The positive and negative

m_z components at the edges stem from the Oersted field. The corresponding power spectral density (PSD), computed from the magnetization dynamics in the central constriction, exhibits no distinct spectral features (panel (i) of [figure 6.6b](#)), indicating the absence of self-oscillations. In contrast, when a superthreshold dc current density $J_{dc}^{\text{above}} = 2.2 \times 10^8 \text{ A/cm}^2$ is applied, persistent excitations localized within the constriction emerge (panel (ii) of [figure 6.6a](#)). The associated PSD shows a broad, enhanced response near the optical eigenmode frequency (panel (ii) of [figure 6.6b](#)), and the oscillations emit short-wavelength ($\leq 50 \text{ nm}$) SWs propagating outward into the surrounding SAF.

When an additional rf current density $J_{rf}^{\text{peak}} = 0.5 \times 10^8 \text{ A/cm}^2$ at 8 GHz is applied simultaneously, part of the self-oscillatory motion becomes phase-locked to the monochromatic drive (panel (iii) of [figure 6.6a](#)). This injection-locking modifies the emitted spin-wave pattern: longer-wavelength magnons consistent with the rf locking frequency dominate the propagation, and the PSD displays a sharp peak at the locking frequency (panel (iii) of [figure 6.6b](#)), confirming synchronization between the dc-driven self-oscillation and the rf excitation. Crucially, this injection locking provides a pathway for detecting dc-induced self-oscillations via SRE, as the synchronized precession produces an oscillatory resistance at the rf frequency. This homodyne detection mechanism requires that the spectrum of the SOT-driven oscillation contains the drive frequency f_1 and that the precession remains phase-locked to the rf current²⁴². It is important to emphasize that the injection-locking simulations are not restricted to the 45 mT field and 8 GHz rf frequency shown here. We performed simulations at external fields of 30, 40, 45, and 50 mT, and for rf frequencies of 4, 5, 6, 7, and 8 GHz. In all cases, injection-locked self-oscillatory dynamics were observed (data not shown).

As shown in panel (ii) of [figure 6.6b](#), in addition to the broad spectrum of self-oscillation modes, the PSD for the OM is the dominant one. At present, the precise origin of the dominance of the optical mode is not fully understood. However, based on our simulations and experimental conditions, we propose the following possible interpretation. The simulations correspond to magnetic fields within approximately $\pm 50 \text{ mT}$, where the magnetization dynamics are primarily governed by the strong antiferromagnetic exchange coupling between the layers. In this low-field regime, the system may naturally favor the excitation of the exchange-dominated optical mode. In contrast, efficient excitation of the acoustic mode generally requires larger external magnetic fields and/or a significant tilting of the

magnetizations away from the antiparallel configuration. However, we emphasize that this explanation remains tentative.

6.3.3 Threshold current and emergence of chiral self-oscillations in SAF

To quantify the magnetoresistance, we applied a small direct current $I_{dc} = 50 \mu A$ through the nano-constriction device and recorded the resulting DC voltage drop using a precision voltmeter while sweeping the magnetic field for two in-plane orientations, $\Phi = 0^\circ$ and 90° , where Φ is the azimuthal angle between the magnetic field ($\mathbf{H}$) and the current direction. Figure 6.7a shows the percentage change in resistance as a function of magnetic field for two in-plane orientations, $\Phi = 0^\circ$ and 90° . The maximum resistance change of approximately 0.2% observed at $\Phi = 90^\circ$ arises because the AMR and GMR contributions add up when the field drives a rotation and relative twisting of the sublattice magnetizations. In contrast, when the field is applied along $\Phi = 0^\circ$, the AMR and GMR contributions cancel each other during the field sweep between the spin-flip and spin-flop fields. This demonstrates that in our SAF the AMR and GMR contributions are of comparable magnitude, each amounting to about 0.1%. Figure 6.7b presents measurements performed at $\Phi = 45^\circ$, where both AMR and GMR contribute simultaneously. In this configuration, sharp resistance variations are observed at fields of approximately ± 50 mT, indicating spin-flop transitions of the SAF. The presence of hysteresis in both the AMR and GMR signals between the spin-flip and spin-flop fields demonstrates the existence of in-plane biaxial anisotropy. This anisotropy most likely originates from the geometric patterning of the nanoconstriction, which defines two easy axes aligned along the constriction edges ($\pm 45^\circ$). From these measurements, the biaxial anisotropy field is estimated to be about 5 mT. Although very weak in strength, this in-plane biaxial anisotropy induced low field spin-flop transition is crucial for the emergence of DC current-driven self-oscillations.

We now present experimental evidence for the emergence of self-oscillations near the spin-flop transition, i.e., at low applied magnetic fields ($|\mu_0 H| < 50$ mT), where no resonant eigenmodes are observed in the absence of a dc current. Figure 6.8a shows the spin rectification voltage V_{SRE} as a function of in-plane magnetic field ($\Phi = 0^\circ$) for three conditions: zero dc bias ($I_{dc} = 0$) and finite dc currents of

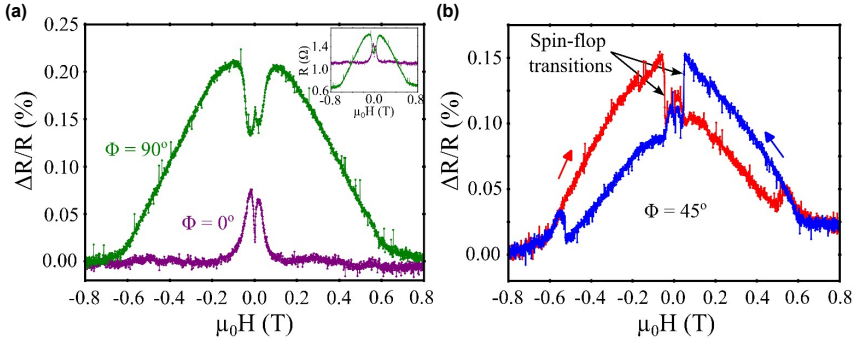


Figure 6.7: AMR and GMR in SAF: (a) Percentage of change in resistance ($\Delta R/R$) vs. the magnetic field H for two different in-plane orientations of H ; $\Phi = 0^\circ$ and 90° for the patterned SAF nanoconstriction. The inset shows the as measured resistance vs. the magnetic field where a constant offset of 445Ω is subtracted from the data. (b) $\Delta R/R$ vs. H for two different sweeping directions and at $\Phi = 45^\circ$.

$I_{dc} = \pm 0.45$ mA, corresponding to a current density of 2.9×10^7 A/cm² in the nanoconstriction. At zero dc current, no SRE features appear at the 8 GHz probing frequency, consistent with the absence of resonant excitation. In contrast, sharp SRE peaks emerge for both positive and negative dc currents, evidencing the onset of resonant excitations. Notably, the sign of these DC current induced SRE peaks depends solely on the polarity of the applied DC current and is independent of the magnetic field direction, in contrast to the conventional SRE behavior observed for linear eigenmodes in the absence of DC bias (section 6.3.1).

We attribute these anomalous threshold current dependent peaks to self-sustained oscillations of Néel order driven by the DC current-induced DL-SOT, which become injection-locked to the fixed rf frequency. The observed current-polarity dependence of the SRE signal provides strong evidence for the chiral nature of the self-oscillations. Importantly, these DC current-induced oscillations differ fundamentally from the linearly polarized acoustic and optical eigenmodes. Near the spin-flop transition, the DL-SOT acts as a staggered effective field on the two antiferromagnetically coupled layers, leading to circularly polarized, chiral auto-oscillatory dynamics whose handedness is selected by the current polarity. Since both AMR and GMR are even under simultaneous reversal of $\mathbf{M}_1$ and $\mathbf{M}_2$, the sign of the SRE voltage is insensitive to the magnetic-field polarity and depends only on the dc current direction. This behavior is consistent with the reversal of

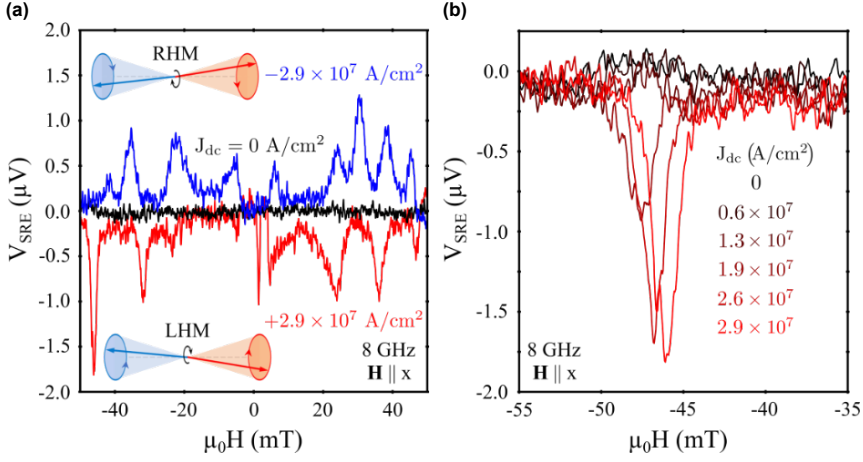


Figure 6.8: (a) Electrical measurements of the spin rectification voltage V_{SRE} associated with self-oscillations, plotted as a function of the applied magnetic field for three dc current densities: 0 (black), -2.9×10^7 A/cm² (blue), and $+2.9 \times 10^7$ A/cm² (red). The rf current is applied at 8 GHz with a power of 15 dBm. Schematics of the two possible chiralities of the self-oscillation modes, left-handed (LHM) and right-handed (RHM), are shown for illustration. (b) Dependence of a representative self-oscillation mode on the dc current density, varied from 0 to 2.9×10^7 A/cm², demonstrating the threshold behavior of the oscillation amplitude. In this measurement, the magnetic field is applied along the bar direction ($\Phi = 0^\circ$).

the spin accumulation at the Pt and Ta interfaces when the current direction is inverted.

These observations agree with theoretical predictions for in-plane antiferromagnetic SHNOs, in which the chirality of the self-oscillations is set by the direction of the applied current²³⁷. Similar electrical detection schemes for magnon mode chirality has been shown for ferrimagnetic systems, where DL torque governs the handedness of ferrimagnetic eigenmodes, with the sign of the rectified voltage reflecting the magnon-mode chirality²⁴⁵. While their work addresses chiral eigenmodes rather than dc-driven self-oscillations, they establish that electrical readout techniques, including SRE, can directly probe magnon chirality in antiferro/ferrimagnetic systems. Using Fourier analysis of the time-dependent AMR and GMR signals from our micromagnetic simulations, which will be shown in the next subsection, we confirm numerically that the polarity of the SRE voltage arising from injection-locked self-oscillations depends solely on the dc current

direction and is independent of the applied magnetic field.

Thus, in the low-field regime near the spin-flop transition, the chirality of the self-oscillation is uniquely controlled by the dc current polarity, an important distinction from ferromagnetic STT and SOT oscillators, where the chirality is always right-handed with respect to the applied field²³⁷.

The threshold behavior of the self-oscillations is shown in [figure 6.8b](#). Focusing on a representative SRE peak within the field window -55 mT to -35 mT, we vary the dc current corresponding to current densities from 0 to 2.9×10^7 A/cm². This resonance appears only above a critical current density of 1.3×10^7 A/cm² and shifts to lower magnetic fields as the current increases. Notably, the critical current density obtained from micromagnetic simulations is approximately one order of magnitude larger than the experimental value. However, we emphasize that such differences between deterministic micromagnetic simulations and experiments are not uncommon in current-driven spin-torque systems. Our simulations do not include thermal fluctuations or Joule-heating-induced temperature rises, both of which can substantially reduce the effective threshold for current-driven dynamical instabilities by assisting the system in overcoming energy barriers and by modifying material parameters such as magnetization, anisotropy, and damping. Temperature-assisted reductions of critical SOT current densities, by up to and beyond one order of magnitude, have been widely reported for SOT-driven switching and related dynamical phenomena. While a quantitative analysis of temperature-assisted threshold reductions for SOT-driven self-oscillations has not yet been reported to our knowledge, the same qualitative trend is physically well-motivated.

In addition to the SRE detection of DC current-driven self-oscillations observed in our SAF with an asymmetric notch-shaped nano-constriction, we have also measured similar self-oscillations in a symmetric nano-constriction with a slightly larger constriction width (approximately 450 nm), as shown in [figure 6.9a](#). [Figure 6.9b](#) presents a colormap of the measured V_{SRE} response as a function of both the applied magnetic field and the DC current density. The threshold current density for the onset of auto-oscillations can be clearly identified and is approximately equal to that observed in the asymmetric device, i.e. $\approx 1.3 \times 10^7$ A/cm². This strongly suggests that the emergence of self-oscillations is a DC current-induced effect, as the rf power was kept constant at 15 dBm for both cases. Here again, the positive (blue) and negative (red) signs of the V_{SRE} peaks are exclusively

observed for negative and positive DC current polarities, respectively, indicating the associated chiralities of the self-oscillation modes.

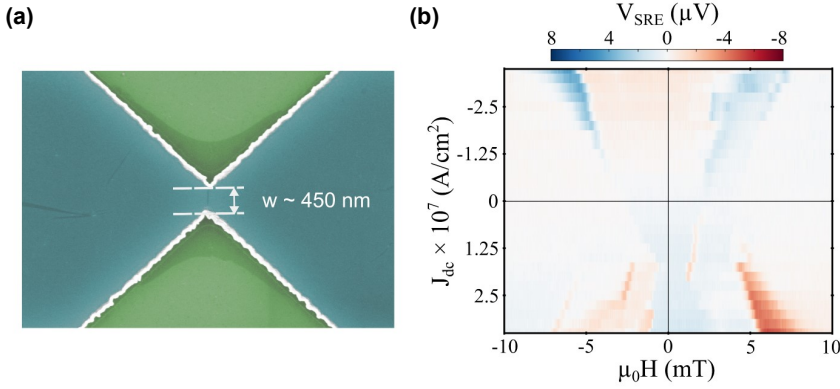


Figure 6.9: (a) Scanning electron microscope image of a symmetric nano-constriction device made from SAF showing the constriction width of 450 nm. (b), Dependence of self-oscillation for the device (a) on the magnetic field and dc current density is shown as a 2D color-plot. In this case a fixed rf probing current with 8 GHz frequency is applied for the injection-locking of the self-oscillations. The magnetic field is oriented along the bar direction.

We observe similar self-oscillation modes when both the DC current and the applied magnetic field are reversed simultaneously; that is, the oscillation characteristics for negative current and negative field closely resemble those for positive current and positive field. Likewise, a comparable symmetry is observed between the cases of positive current with negative field and negative current with positive field. This behavior reflects the inherent symmetry of DL-SOT with respect to current and field reversal. The self-oscillation SRE peak at 8 GHz shifts only slightly with increasing DC current density, showing a magnetic field variation of approximately 1 mT. This minor shift can be attributed to a DC current-induced FL torque, which modifies the total effective magnetic field and consequently causes a slight change in the resonance condition.

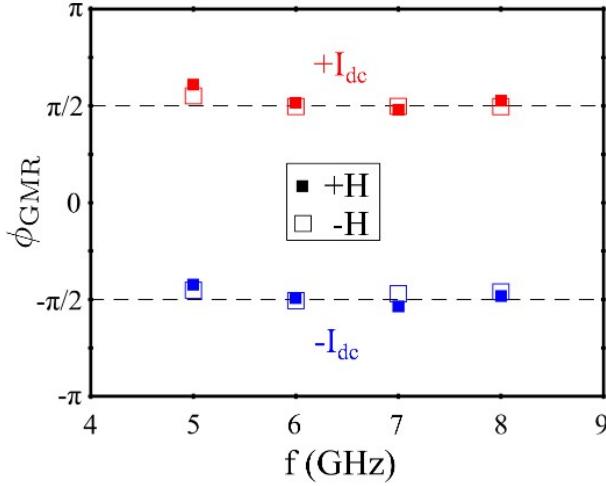


Figure 6.10: The phase of the GMR response obtained from FFT at a given frequency is plotted for four different injection-locking rf frequencies. The squares and circles correspond to positive and negative magnetic polarities respectively with a magnitude 45 mT. The red and blue color corresponds to positive and negative polarity of the direct current. The micromagnetic simulation is carried out for $|I_{\text{dc}}| = 1.9 \text{ mA} > I_{\text{TH}}$ and rf current with 0.5 mA peak magnitude for the respective frequencies.

6.3.4 Chirality of the direct current induced self oscillations in SAF from micromagnetic simulation

Given that the chiralities of the DL torque induced precession are opposite for positive and negative current polarities, the observed sign change in the DC current-driven SRE signal is solely attributed to the reversal of the direct current polarity. To further substantiate this behavior, we have conducted a detailed analysis of the phase relationships in the rectified AMR and GMR responses, extracted via FFT at specific frequencies. This analysis was performed for both opposing magnetic field and current polarities during micromagnetic simulations of injection-locked self-oscillations in our SAF system. By examining the phase shifts, we can gain insight into how the chirality of the precession depends on the polarities of the applied magnetic field or current. The GMR response in this case depends on the relative magnetization alignment of the two ferromagnetic Py layers in the SAF, which is described by the tilt angle ($\theta_{\text{tilt}} = \cos^{-1}(\vec{M}_1 \cdot \vec{M}_2 / (|\vec{M}_1| |\vec{M}_2|))$). GMR

response is then proportional to $\sin^2(\theta_{\text{tilt}}/2)$. On the other hand, the AMR is influenced by the relative orientation of the magnetizations of the two ferromagnetic Py layers with respect to the direction of current flow. The AMR is proportional to $\sin^2(\cos^{-1}(M_{x_i}/|\vec{M}_i|))$ where $i = 1, 2$ denotes the two ferromagnetic Py layers. To estimate the total AMR contribution from both layers, a parallel resistor model is used. Fig. [Figure 6.10](#) presents the phase of the rectified GMR response, obtained via FFT for four different injection-locking rf frequencies and for opposite polarities of both the magnetic field and current. It is apparent that for all rf input frequencies, the phase of the rectified GMR response changes by π when the current polarity is reversed. As expected, no phase reversal is observed when the magnetic field polarity is reversed. The phase analysis from the rectified AMR response shows similar results however for certain frequencies the AMR rectification is absent, hence is not shown here.

6.3.5 Potential signatures of SOT-induced chaotic dynamics near the spin-flop transition

To investigate the rich nonlinear dynamics predicted for SOT-driven antiferromagnetic oscillators near the spin-flop transition, we performed spin-rectification measurements over a frequency range of 5-8 GHz and compared the behavior of the acoustic eigenmode ([section 6.3.1](#)) with the DC-driven self-oscillations presented in [section 6.3.3](#). In the absence of a DC bias, only the acoustic eigenmode is detected, exhibiting a resonance field that shifts monotonically to higher values with increasing RF frequency ([figure 6.11a](#) and [figure 6.11c](#)), consistent with the analytical model of Section 1.

When the DC current density exceeds the critical threshold $J_{\text{dc}}^{\text{above}} = 2.5 \times 10^7 \text{ A/cm}^2$, a sequence of additional sharp rectification peaks emerges near zero magnetic field, i.e., in close proximity to the spin-flop transition ([figure 6.11d](#)). These peaks, absent under zero DC bias ([figure 6.11c](#)), correspond to the injection-locked self-oscillations discussed in [section 6.3.3](#). Unlike the acoustic mode ([figure 6.11b](#)), the field positions of these low-field peaks remain essentially invariant with rf frequency, indicating a qualitatively different dynamical regime. We attribute their frequency independence to the broadband PSD of the underlying DC-driven self-oscillations, in agreement with the micromagnetic simulation results of [section 6.3.2](#) (cf. [figure 6.6](#)).

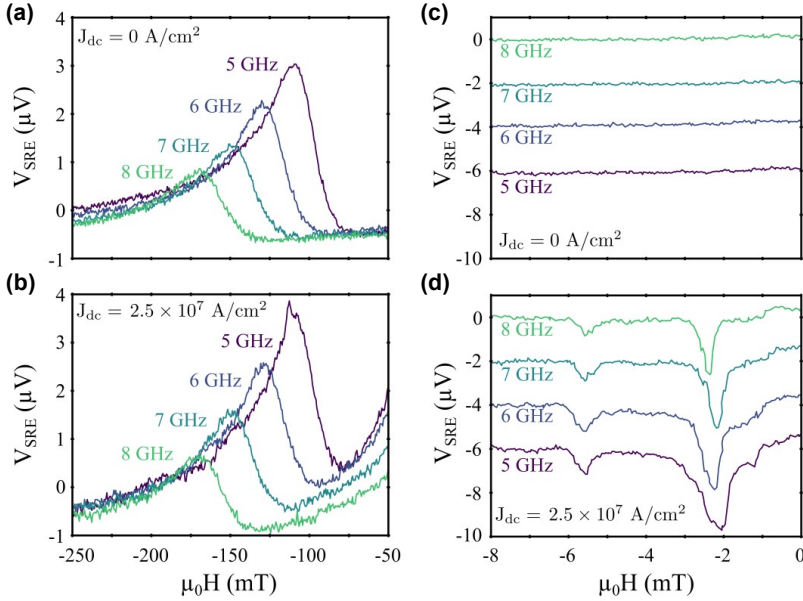


Figure 6.11: (a) and (b) display the acoustic eigenmode measured via the SRE for rf frequencies between 5 and 8 GHz under $J_{dc} = 0$ and $J_{dc} = +2.9 \times 10^7$ A/cm², respectively. Panels (c) and (d) show the corresponding SRE measurements in the low-field regime where injection-locked self-oscillations emerge, again for rf frequencies from 5 to 8 GHz and for $J_{dc} = 0$ and $J_{dc} = +2.9 \times 10^7$ A/cm², respectively. For clarity, the SRE traces in (c) and (d) are vertically offset to improve visualization of the individual self-oscillation peaks. In all measurements, the magnetic field is applied nominally along the bar direction with a slight intentional misalignment, and the rf power is fixed at 15 dBm for all frequencies.

To further elucidate the origin of these excitations, we performed macrospin simulations using experimentally determined RKKY exchange and anisotropy parameters. Whereas micromagnetic simulations capture the spatial structure of the dynamics, the macrospin model provides access to extended dynamical evolution over long timescales. These simulations (figure 6.12) reveal complex nonlinear behavior near the spin-flop transition. In particular, Lyapunov exponent analysis and bifurcation diagrams indicate transitions between stable and chaotic oscillatory regimes, depending on the applied DC current and magnetic field. These results are consistent with theoretical predictions²³⁹ and suggest that the sudden appearance of multiple threshold-dependent resonances, combined with their insensitivity to rf frequency, constitutes a possible experimental signature of

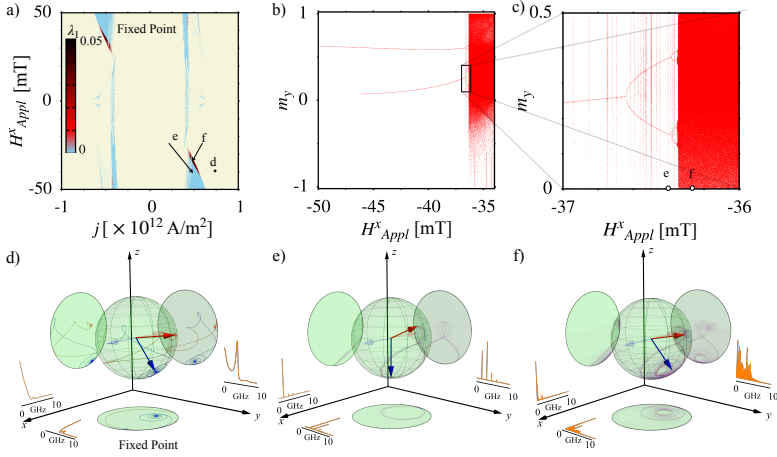


Figure 6.12: (a) Color density plot of the maximum Lyapunov exponent λ_1 as a function of the applied field along x -axis H_{ext}^x in mT and the current density j in A/m^2 . This plot illustrates regions of chaotic behavior ($\lambda_1 > 0$) and stable dynamics ($\lambda_1 = 0$). (b) Bifurcation diagram for a current density $j = 0.5 \times 10^{12} \text{ A/m}^2$ as a function of H_{ext}^x . This diagram shows the transition from periodic to chaotic behavior as the field increases. (c) Enlarged view of the bifurcation diagram highlighting a period-doubling route to chaos. (d-f) Different dynamic states of the system: (d) fixed point state indicating a stable equilibrium; (e) periodic dynamic state marked as (e) in the bifurcation diagram, showing the transient period before reaching the final periodic state; and (f) chaotic dynamic state marked as (f) in the bifurcation diagram, illustrating the complex and unpredictable behavior typical of chaotic systems. The associated FFTs are shown for each component: x in the xy -plane, z in the xz -plane and y in the yz -plane.

chaotic SOT-driven dynamics in the SAF system. Nevertheless, direct experimental discrimination from broadband nonlinear dynamics requires future noise or phase-resolved measurements.

6.4 Conclusion

We have demonstrated a SOT-driven antiferromagnetic oscillator based on a SAF nanoconstriction, combining electrical detection with micromagnetic and macrospin simulations. Spin-rectification spectroscopy revealed magnetic-field-dependent excitations of the linearly polarized acoustic and optical eigenmodes and, critically, the emergence of additional resonances that appear only when the

applied DC current exceeds a threshold.

These additional resonances are consistent with self-sustained oscillations of the antiferromagnetic order, which are absent without the DC drive. When a rf probe current is superimposed, the self-oscillations become injection-locked to the rf frequency and manifest as sharp, field-resonant spin-rectification peaks. The polarity of the rectified voltage depends exclusively on the direction of the DC current, confirming that the oscillation chirality is set by the current polarity, a key distinction from ferromagnetic spin-torque oscillators.

Micromagnetic simulations incorporating experimentally determined material parameters and realistic device geometry reproduce the emergence of these chiral self-oscillations and their injection locking. Complementary macrospin simulations, supported by Lyapunov-exponent and bifurcation analyses, reveal that SAF-based oscillators exhibit both stable and chaotic dynamical regimes near the spin-flop transition. These numerical results support the interpretation that the frequency-insensitive, multi-peak spin-rectification response observed above the threshold current may signal the onset of chaotic behavior, a theoretically predicted but previously unobserved feature of antiferromagnetic spintronic oscillators. We emphasize that the present work does not provide direct spectroscopic detection of self-oscillations or chaotic dynamics, but instead establishes electrical signatures and operating regimes consistent with such behavior.

Together, these results establish synthetic antiferromagnets as a powerful and experimentally accessible platform for exploring chiral and non-linear antiferromagnet-like spin dynamics, with relevant for future ultrafast real antiferromagnetic systems. The combination of tunable nonlinear behavior, broadband operation, and inherent stochasticity positions SAF-based oscillators as promising building blocks for emerging applications in neuromorphic computing, signal processing, and unconventional information encoding.

VII

Outlook and future perspectives

7.1 Overall summary of the thesis

This thesis establishes SAFs as a versatile and experimentally accessible platform for SOT-driven antiferromagnetic spintronics, spanning static non-volatile memory functionality, dynamic real-space imaging, and nonlinear coherent oscillations. First, we demonstrated a room-temperature, nonvolatile analog memristor based on a magnetically compensated out-of-plane SAF with PMA. Information is encoded in the ratio of reversed antiferromagnetic domains and written via DL SOT-driven DW motion. The absence of stray fields ensures remarkable reproducibility and stability of the analog state. Nonlinear magnetotransport aided by DL SOT provides electrical readout, establishing SAFs as artificial synapses exhibiting pulse-amplitude-dependent plasticity: a distinct neuromorphic functionality rooted in antiferromagnetic domain control.

Second, we extended this control into the dynamic regime by developing a SOT-enabled polar tr-MOKE approach capable of imaging in-plane compensated

SAF domains in real space. By exploiting the staggered nature of the DL spin-orbit fields, we established a robust optical contrast mechanism for detecting the reversed Néel order. Phase-resolved tr-MOKE measurements revealed a 180° phase shift of domain oscillations across DWs, directly encoding the Néel vector orientation. We further demonstrated reversible DC-current-driven DW motion and identified a geometry-confined localized excitation mode within the microconstriction.

Third, with the help of electrical SRE measurements, we demonstrated a SOT-driven SAF nano-oscillator exhibiting field-dependent acoustic and optical eigenmodes as well as threshold-induced chiral self-oscillations. The chirality of these auto-oscillating modes is controlled by DC current polarity and these modes can injection-lock to external microwave signals. Micromagnetic and macrospin simulations reveal both stable and chaotic dynamical regimes near the spin-flop transition, positioning SAFs as a platform for nonlinear and potentially stochastic spintronic functionality.

Taken together, these results show that SOT in compensated SAFs enable: deterministic, nonvolatile domain-state control, direct real-space access to dynamic Néel order of an in-plane SAF and coherent, chiral, and nonlinear oscillatory behaviors. This unified framework demonstrates that SAFs are not merely model systems for AFs, but powerful device platforms for functional spintronics. In a broader context, this thesis responds directly to the longstanding challenge articulated in antiferromagnetic spintronics: how to electrically control, dynamically probe, and functionally exploit compensated magnetic order. By employing SAFs as tunable and experimentally accessible model systems, we bridge the conceptual advantages of real antiferromagnets. Looking ahead, the concepts established in this thesis extend well beyond SAFs as model systems. The ability to electrically write, dynamically probe, and nonlinearly drive compensated spin textures provides a transferable framework for emerging classes of materials, including crystalline antiferromagnets and recently identified altermagnets with unconventional symmetry, ultrafast THz dynamics and anomalous transport properties. At the same time, the convergence of analog memory, geometry-engineered spin dynamics, chirality control, and nonlinear oscillatory behavior suggests a natural pathway toward hardware implementations of neuromorphic and probabilistic computing.

Building on the findings of this thesis, the following section outlines several

forward-looking directions which are currently being pursued, including direct spectral and optical detection of SOT-driven self-oscillations, current-controlled frequency tuning in exchange-dominated regimes of self-oscillation, nonlinear magnonic phenomena such as frequency-comb generation, and ultrafast nano-scale imaging of SOT-driven spin dynamics by utilizing time-resolved Lorentz microscopy.

7.2 Future perspectives

7.2.1 Spectrum-analyzer and Brillouin light scattering detection of SOT-driven self-oscillations

In [chapter VI](#), we employed a single experimental approach, the SOT-driven SRE technique, to investigate both the intrinsic antiferromagnetic acoustic and optical eigenmodes of the easy-plane in-plane SAF and the emergence of additional dc-current-driven self-oscillatory modes. While this method provides compelling evidence for current-induced dynamical phenomena, a conclusive demonstration of dc-current-induced self-oscillations and potential chaotic dynamics requires complementary spectroscopic characterization. In particular, direct measurements of the spectral emission of the SOT-driven SAF nano-oscillator are essential, both electrically, using a high-frequency spectrum analyzer, and optically via Brillouin light scattering (BLS) spectroscopy. Micro-focused BLS further offers the capability of spatially resolved, scanning-based microscopy, enabling direct probing of the spatial confinement of dc-current-driven self-oscillations in SAF structures without relying on injection-locking schemes. We have initiated systematic BLS investigations, and the first measurements already reveal promising signatures of current-induced modulation of the antiferromagnetic eigenmodes.

As shown in [figure 7.1a](#), a magnetic field of +250, mT applied perpendicular to the bar device used in [chapter V](#) (i.e., at $\alpha = 90^\circ$) excites the acoustic eigenmode at approximately 16GHz in the absence of dc current ($I_{dc} = 0\text{mA}$). Upon application of a positive dc current of +15mA, a pronounced enhancement of both the Stokes and anti-Stokes intensities in the BLS spectrum is observed, indicating the anti-damping character of the current-induced SOT. Conversely, reversing the current polarity to -15mA results in a substantial suppression of the acoustic mode

amplitude, consistent with the damping action of the SOT.

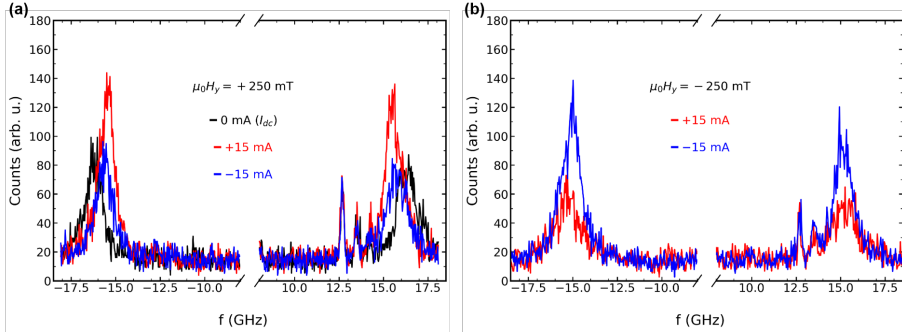


Figure 7.1: Micro-focused BLS spectra of the acoustic eigenmode in the SAF constriction device (the one used in [chapter V](#)) under perpendicular magnetic fields. (a) Stokes and anti-Stokes spectra measured at +250 mT ($\alpha = 90^\circ$) for three dc currents: 0 mA (black), +15 mA (red), and -15 mA (blue). The spectra show the enhancement of the mode for positive current (anti-damping) and suppression for negative current (damping) relative to zero current. (b) Spectra under reversed magnetic field -250 mT, measured for +15 mA (red) and -15 mA (blue). Reversal of the field polarity inverts the current-induced damping and anti-damping behavior, consistent with the symmetry expected for the DL SOT in the SAF. Some small artifact peaks are visible in both panels; these do not change with current and can be ignored.

The symmetry of the damping and anti-damping response is further corroborated by reversing the magnetic field polarity to -250 mT. In this configuration, -15 mA produces an anti-damping effect, whereas +15 mA induces damping of the acoustic eigenmode, in contrast to the +250 mT case. This field-current symmetry is characteristic of a DL SOT and confirms the presence of a current-induced DL SOT component in our SAF system capable of controllably modifying the eigenmode spectrum. These initial BLS measurements, although not yet fully optimized, establish a clear proof-of-principle. Further systematic studies are planned to sensitively detect dc-current-induced self-oscillations, particularly within the spin-flop transition regime discussed in [chapter VI](#).

In parallel, we intend to perform direct electrical spectral emission measurements of the SAF nano-oscillators using a high-frequency spectrum analyzer. At present, the signal-to-noise ratio of the electrical detection is limited by the relatively small AMR and GMR responses, which yield relative resistance variations on the order of $\sim 0.1\%$. This limitation can be mitigated through established

strategies for enhancing microwave output power and detection sensitivity. These include high-contrast electrical readout via MTJ-based tunneling magnetoresistance or optimized spin-valve geometries, impedance-matched microwave circuitry, and mutual synchronization of multiple nano-oscillators. Such approaches are routinely employed in ferromagnetic STT and SOT nano-oscillators and are directly transferable to SAF-based devices.

The combination of electrical spectrum analysis and optical BLS spectroscopy will ultimately enable comprehensive noise spectroscopy of the SAF nano-oscillator. This will provide critical insight into linewidth broadening, phase noise characteristics, and potential signatures of chaotic dynamics, thereby allowing a rigorous verification of the SOT-driven nonlinear dynamical states discussed in this thesis.

7.2.2 Current-controlled frequency tunability in exchange-dominated self-oscillations

Beyond the chirality control via current polarity in SAF-based nano-oscillators described in [chapter VI](#), and in addition to their chaotic dynamics enabling broadband excitation of self-oscillatory modes, SAF systems offer a further key advantage: the possibility of achieving an almost linear frequency tunability with dc current magnitude. If experimentally realized, this would constitute a significant improvement over conventional FM-based oscillators, in which current-driven frequency tunability is typically weak and often requires additional voltage-controlled magnetic anisotropy (VCMA) mechanisms to achieve substantial frequency modulation.

Preliminary macrospin simulations indicate that reducing the effective in-plane anisotropy, feasible via interface engineering, may enable exchange-dominated self-oscillations exhibiting cleaner, more linear frequency tunability and a reduced tendency toward chaotic behavior. The physical mechanism underlying these linear, exchange-dominated SOT-driven oscillations is schematically illustrated in [figure 7.2a](#), [figure 7.2b](#), and [figure 7.2c](#). In this regime, a strong external magnetic field ($\mathbf{H}$), applied close to the spin-flip field, cants the magnetizations of the two FM layers toward the field direction. This canting generates an exchange field ($\mathbf{H}_{ex}$), which, together with the current-induced DL SOT, drives

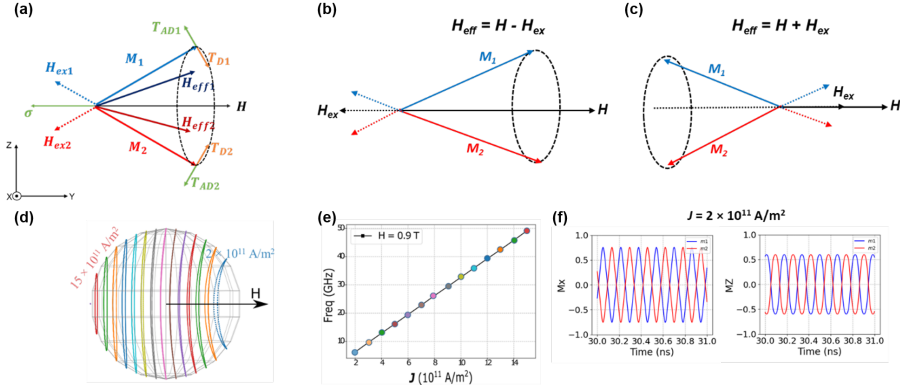


Figure 7.2: Schematic illustration and macrospin simulation results of exchange-dominated, linearly frequency-tunable SOT-driven self-oscillations in a low-anisotropy SAF system. (a) Schematic configuration under a strong external magnetic field $\mathbf{H}$ applied close to the spin-flip field, canting the magnetizations of the two FM layers. The resulting exchange field $\mathbf{H}_{ex}$, together with the damping-like SOT, drives sustained self-oscillations. (b) Low-current regime: the magnetizations remain predominantly tilted toward the applied field direction, such that the exchange field partially subtracts from $\mathbf{H}$, yielding an effective field $\mathbf{H}_{eff} = \mathbf{H} - \mathbf{H}_{ex}$. (c) High-current regime: for sufficiently strong anti-damping SOT, the magnetizations cant toward the spin-polarization direction (opposite to $\mathbf{H}$), causing the exchange field to align parallel to the applied field and add constructively, $\mathbf{H}_{eff} = \mathbf{H} + \mathbf{H}_{ex}$. (d) Macrospin simulation showing the evolution of the precessional orbits with increasing current amplitude; the cone angle increases approximately linearly in the first half of the Bloch sphere and decreases in the second half as the system transitions between the two canting regimes. (e) Corresponding oscillation frequency as a function of current amplitude, demonstrating an approximately linear tunability over a broad range. (f) Time evolution of the x - and z -components of the magnetization illustrating the steady-state self-oscillatory dynamics. A variation of current amplitude by approximately one order of magnitude enables frequency tuning from $\sim 10\text{GHz}$ to $\sim 50\text{GHz}$.

sustained self-oscillations. The magnetizations precess around the effective field $\mathbf{H}_{eff} = \mathbf{H} \pm \mathbf{H}_{ex}$, where the $\pm$ sign depends on the relative orientation of the two FM layer magnetizations, as depicted in [figure 7.2b](#) and [figure 7.2c](#) with respect to $\mathbf{H}$. Depending on the configuration, the exchange field either subtracts from or adds to the applied field. For low current amplitudes, the magnetizations remain predominantly tilted toward the applied magnetic field during self-oscillation. As the current increases, the strength of the DL-SOT grows, leading to a larger precessional cone angle. Consequently, the effective exchange field is progressively reduced, meaning that a smaller portion of $\mathbf{H}_{ex}$ is subtracted from $\mathbf{H}$. This results

in an increase of the net effective field $\mathbf{H}_{eff}$. Since the oscillation frequency is determined by the magnitude of $\mathbf{H}_{eff}$, the frequency increases approximately linearly with current.

For sufficiently strong anti-damping SOT, the magnetizations become canted toward the spin-polarization direction, which is opposite to the applied magnetic field direction (see [figure 7.2c](#)). In this configuration, the exchange field becomes parallel to the applied magnetic field and therefore adds constructively to it. With further increase in current amplitude, the magnetizations are canted more strongly toward the spin-polarization direction, enhancing the contribution of $\mathbf{H}_{ex}$ to $\mathbf{H}_{eff}$. As a result, the total effective field increases further, leading to a continued increase in oscillation frequency.

The macrospin simulation results for a SAF system with reduced anisotropy are presented in [figure 7.2d](#). With increasing current amplitude, the precessional orbits evolve such that the cone angle increases approximately linearly in the first half of the Bloch sphere and decreases in the second half, reflecting the transition between the two canting regimes described above. The corresponding linear dependence of oscillation frequency on current amplitude is shown in [figure 7.2e](#), while the time evolution of the x - and z -components of the magnetization is presented in [figure 7.2f](#). Notably, a change in current amplitude by approximately one order of magnitude enables frequency tuning from roughly 10GHz to 50GHz, highlighting the exceptional tunability potential of exchange-dominated SAF oscillators.

7.2.3 Magnon frequency comb in SAFs

A frequency comb is a spectrum which consists of a series of discrete equidistant spectral lines in the frequency domain. It was originally proposed and realized in optical systems with immense importance in optical atomic clock, light detection and ranging (LiDAR), coherent optical communication, astronomy and cosmology etc. which led to the Nobel Prize in Physics in 2005^{246,247}. Despite the rapid development in photonic frequency combs, there has been only a limited studies on the magnetic counterpart i.e. magnon frequency comb (MFC). Similar to photons, magnons are bosonic quasiparticles and should exhibit similar phenomena. There have been theoretical predictions on MFC arising from non-linear

magnon-skyrmion scattering and interaction between propagating spin wave and oscillating DW^{59,248}. MFC has been realized experimentally in a magnetomechanical system owing to the magnon-phonon coupling and magnetostriction effect²⁴⁹. Here we show the evidence of the existence of a MFC in a SAF by performing micromagnetic simulations. The inherent nonlinearities of the magnetization dynamics described by LLG equation including the nonlinear current induced DL SOT and interlayer exchange torque can lead to parametric amplifications and three and four magnon interactions at sufficient input RF power²⁵⁰. The three and four magnon scattering processes are shown in figure 7.3. This nonlinear cascaded magnon-magnon interactions will lead to generation of series of distinct frequencies similar to the optical frequency comb.

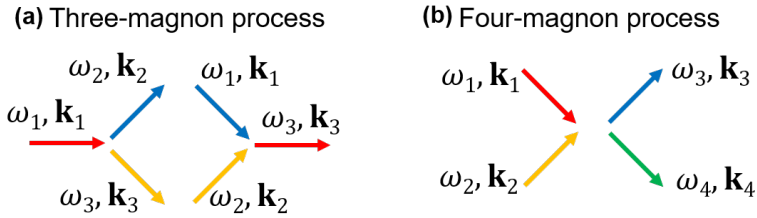


Figure 7.3: Schematic illustration of three and four magnon scattering processes. Adapted from [250].

Figure 7.4a and figure 7.4b shows the results of fast Fourier transform (FFT) performed on the time evolution of the net magnetization taken from the 20 nm x 20 nm central region within the nano-constriction of the SAF (same nano-constriction used for simulations in chapter VI). A small magnetic field of strength +45 mT is applied slightly misaligned to the bar orientation. In this case, the dc current is kept below the threshold for the self-oscillation whereas the rf current strength is gradually increased at a fixed frequency of 8 GHz. At 1.0 mA peak rf amplitude, as shown in Figure figure 7.4a, MFC effect can be observed with a mode spacing of 4 GHz. This might correspond to splitting of the incident magnon ($\omega_p = 8$ GHz) into two different magnons with opposite wave vectors and frequency $\omega_k = \omega_p/2 = 4$ GHz, which subsequently interact (annihilation and confluence) in a cascaded manner to produce sharp modes every 4 GHz interval. Again, as shown in figure 7.4b, secondary MFC effect with finer mode spacing (0.25 GHz) can be observed around all the 4 GHz spaced MFC modes. This shows, SAF along with the complex antiferromagnetic dynamics at the critical spin-flop transition

region can also exhibit rich MFC phenomena as well. The MFC effect in SAF can be detected by using frequency resolved BLS spectroscopy, spectrum analyzer, SRE and by trMOKE measurements. Further theoretical and experimental investigations of magnetization dynamics and nonlinear magnon–magnon interactions in SAFs, as well as in real AF materials, are being pursued in our group. Particular emphasis will be placed on developing a comprehensive understanding of the underlying nonlinear coupling mechanisms and their impact on coherent and incoherent dynamical regimes.

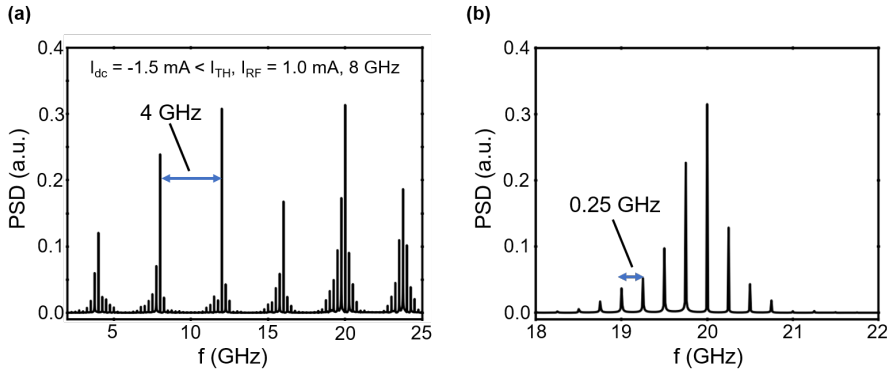


Figure 7.4: MFC in SAF nano-constriction: a) FFT spectrum at $I_{dc} = -1.5$ mA ($< I_{TH}$, threshold current for self-oscillations) and $I_{rf} = 1.0$ mA peak amplitude at 8 GHz showing MFC effect with mode spacing at 4 GHz interval. b) magnified FFT spectrum around the 20 GHz line from a) showing secondary MFC with mode spacing of 0.25 GHz

7.2.4 Ultrafast Lorentz microscopy of SOT-driven dynamics in SAFs

As demonstrated in [section 6.3.2](#) by micromagnetic simulations, non-reciprocal magnons are excited when the SAF nano-constriction is driven by a current-induced SOT. These antiferromagnetic spin waves possess very short wavelengths, typically below 50nm, and propagate away from the nano-constriction under dc current excitation. Resolving such short-wavelength excitations exceeds the spatial resolution limits of optical techniques such as scanning tr-MOKE and micro-focused BLS, which are fundamentally restricted to approximately sub-

micron even when employing high-numerical-aperture objectives. Consequently, direct optical imaging of these propagating antiferromagnetic magnons is not feasible.

To overcome this limitation, time-resolved real-space Lorentz microscopy imaging of these excitations will be pursued in collaboration with our collaborators. This technique offers substantially enhanced spatial resolution and is therefore well suited to resolve sub-100-nm spin-wave modes. In particular, the group of Prof. Schäfer at the University of Regensburg is currently developing an ultrafast transmission electron microscopy (TEM) platform that enables in situ excitation of SAF nano-constrictions by rf currents (up to 20GHz) as well as dc currents. For these experiments, the SAF thin films employed in [chapter VI](#) will be deposited onto free-standing silicon membranes to ensure electron transparency.

For time-resolved imaging of the oscillatory SAF spin waves, an ultrafast Lorentz microscopy approach will be employed, which is sensitive to the instantaneous in-plane magnetization. The temporal resolution is determined by the electron pulse duration (on the order of a few hundred femtoseconds) and is therefore significantly shorter than the characteristic timescale of spin precession. In contrast to SRE-based measurements in SAF systems, time-resolved Lorentz microscopy provides access to the full projected in-plane magnetization distribution. This enables simultaneous visualization of localized magnon modes within the driven nano-constriction as well as the emitted propagating spin waves, thereby offering comprehensive real-space insight into the nonlinear magnon dynamics of SAF nano-oscillators.

A

Rectification for the acoustic and optical eigenmodes in an in-plane SAF

A.1 AMR rectification for acoustic mode

We consider an easy-plane in-plane SAF system under the following physical conditions:

- **Current Flow:** The charge current flows along the $\hat{x}$ axis.
- **Spin Polarization:** Top Pt and bottom Ta layer generates via SHE a spin accumulation with polarization $\sigma \parallel \hat{y}$ which is injected into both the ferromagnetic layers of the SAF.
- **Magnetic Configuration:** An external field $\mathbf{H}$ is applied nominally along $\hat{x}$ with strength less than that of the spin-flip field. The magnetizations $\mathbf{M}_1$ and $\mathbf{M}_2$ are canted toward the field. We will consider both cases, field

colinear with the current direction ($\mathbf{H} \parallel \mathbf{J} \parallel \hat{x}$) and a misaligned direction with respect to the current direction by an angle ψ .

- **Anisotropy:** Strong easy-plane in-plane anisotropy is considered such that the out-of plane component of the magnetization is negligible ($m_z \approx 0$).

At equilibrium, the normalized magnetization vectors are:

$$\mathbf{m}_{1,0} = (\cos \theta, \sin \theta), \quad (\text{A.1})$$

$$\mathbf{m}_{2,0} = (\cos \theta, -\sin \theta), \quad (\text{A.2})$$

where θ represents the canting angle relative to the $\hat{x}$ -axis. The total resistance of the SAF is modeled as two resistors in parallel, representing the independent contributions of each magnetic layer:

$$R_{\text{eq}}(t) = \frac{R_1(t)R_2(t)}{R_1(t) + R_2(t)}. \quad (\text{A.3})$$

The AMR for each layer i is given by:

$$R_i(t) = R_0 + \Delta R (\mathbf{m}_i \cdot \hat{x})^2 = R_0 + \Delta R m_{x,i}^2. \quad (\text{A.4})$$

To find the dynamical resistance change, we Taylor expand R_{eq} to first order (for small changes) in δR_i :

$$\delta R_{\text{eq}} \approx \frac{\partial R_{\text{eq}}}{\partial R_1} \delta R_1 + \frac{\partial R_{\text{eq}}}{\partial R_2} \delta R_2. \quad (\text{A.5})$$

Assuming symmetric layers ($R_1 \approx R_2 \approx R$), this simplifies to:

$$\delta R_{\text{eq}} = \frac{1}{4} (\delta R_1 + \delta R_2). \quad (\text{A.6})$$

Now, the goal is to estimate δR_i and substitute it in the above equation to find the change of the total resistance (δR_{eq}). From [equation \(A.4\)](#),

$$\delta R_i = 2\Delta R m_{x,i} \delta m_{x,i}. \quad (\text{A.7})$$

Considering a small deviation δm ($\delta m \ll 1$), then Taylor expanding and keeping only the first order term and subsequently considering the conservation

of the magnetization magnitude ($|\mathbf{m}| = 1$), the equilibrium satisfies the condition:

$$m_x^2 + m_y^2 + m_z^2 = 1. \quad (\text{A.8})$$

Thus,

$$2(m_x \delta m_x + m_y \delta m_y + m_z \delta m_z) = 0. \quad (\text{A.9})$$

As the z component is negligible due to the strong in-plane anisotropy,

$$m_x \delta m_x + m_y \delta m_y = 0 \Rightarrow \delta m_x = -\frac{m_y}{m_x} \delta m_y. \quad (\text{A.10})$$

Applying this to each layer:

- **Layer 1:** $\delta m_{x,1} = -\frac{\sin \theta}{\cos \theta} \delta m_y = -\tan \theta \delta m_y.$
- **Layer 2:** $\delta m_{x,2} = -\frac{-\sin \theta}{\cos \theta} \delta m_y = +\tan \theta \delta m_y.$

The first-order variation in resistance for a single layer is $\delta R_i = 2\Delta R m_{x,i} \delta m_{x,i}$. Therefore:

$$\delta R_1 = 2\Delta R (\cos \theta) (-\tan \theta \delta m_y) = -2\Delta R \sin \theta \delta m_y, \quad (\text{A.11})$$

$$\delta R_2 = 2\Delta R (\cos \theta) (+\tan \theta \delta m_y) = +2\Delta R \sin \theta \delta m_y. \quad (\text{A.12})$$

Summing these for the total equivalent resistance from [equation \(A.6\)](#):

$$\delta R_{\text{eq}} = \frac{1}{4} (-2\Delta R \sin \theta \delta m_y + 2\Delta R \sin \theta \delta m_y) = 0. \quad (\text{A.13})$$

Conclusion: For a perfectly symmetric SAF with the magnetic field applied collinear to the current direction, the AMR rectification of the acoustic mode vanishes. A non-zero spin-rectification voltage in this mode requires a breaking of symmetry between the two layers either by having two unequal FM layers or by applying a misaligned field. The later option is what we will focus next.

Now, misaligning the field by ψ , the equilibrium magnetizations are no longer symmetric:

$$\mathbf{m}_{1,0} = (\cos \theta_1, \sin \theta_1), \quad (\text{A.14})$$

$$\mathbf{m}_{2,0} = (\cos \theta_2, -\sin \theta_2), \quad (\text{A.15})$$

where $\theta_1 \neq \theta_2$. The first order variation in the resistance of the individual layers become from [equation \(A.4\)](#):

$$\delta R_1 = -2\Delta R \sin \theta_1 \delta m_y, \quad (\text{A.16})$$

$$\delta R_2 = +2\Delta R \sin \theta_2 \delta m_y. \quad (\text{A.17})$$

Subsequently, the change in the equivalent resistance from [equation \(A.6\)](#) becomes,

$$\delta R_{\text{eq}} = \frac{\Delta R}{2} (\sin \theta_2 - \sin \theta_1) \delta m_y. \quad (\text{A.18})$$

Note, here if $\theta_1 = \theta_2$, then $\delta R_{\text{eq}} = 0$. Also, the driving torque amplitude is controlled by the projection of the magnetization relative to the fixed spin polarization. Thus,

$$\tau_{\text{DL}} \propto \delta m_y \propto \cos \psi. \quad (\text{A.19})$$

Hence, when $\psi \rightarrow 90^\circ$, $\delta R_{\text{eq}} \rightarrow 0$. The applied radio-frequency current is taken as

$$I(t) = I_{\text{rf}} \cos(\omega t). \quad (\text{A.20})$$

The dc rectified voltage is defined as the time average of the product of the current and the time-dependent resistance,

$$V_{\text{SRE}} = \langle I(t) \delta R_{\text{eq}}(t) \rangle. \quad (\text{A.21})$$

Substituting [equation \(A.18\)](#) and [equation \(A.20\)](#) gives

$$V_{\text{SRE}} = \frac{\Delta R}{2} (\sin \theta_2 - \sin \theta_1) \langle I_{\text{rf}} \cos(\omega t) \delta m_y(t) \rangle. \quad (\text{A.22})$$

Up to this point, no dynamical assumptions have been made. For small-angle magnetization precession driven at angular frequency ω , the transverse magnetization

component can be written as

$$\delta m_y(t) = \Re[\tilde{m}_y(\omega) e^{i\omega t}], \quad (\text{A.23})$$

where $\tilde{m}_y(\omega)$ is the complex response amplitude of the acoustic mode. Now, substituting [equation \(A.23\)](#) into [equation \(A.22\)](#) yields

$$V_{\text{SRE}} = \frac{\Delta R}{2} (\sin \theta_2 - \sin \theta_1) I_{\text{rf}} \langle \cos(\omega t) \Re[\tilde{m}_y e^{i\omega t}] \rangle. \quad (\text{A.24})$$

Using the identity

$$\Re[\tilde{m}_y e^{i\omega t}] = \frac{1}{2} (\tilde{m}_y e^{i\omega t} + \tilde{m}_y^* e^{-i\omega t}), \quad (\text{A.25})$$

and performing the time average, only the dc term survives,

$$\langle \cos(\omega t) \Re[\tilde{m}_y e^{i\omega t}] \rangle = \frac{1}{2} \Re[\tilde{m}_y]. \quad (\text{A.26})$$

Therefore, the rectified voltage becomes

$$V_{\text{SRE}} = \frac{\Delta R I_{\text{rf}}}{4} (\sin \theta_2 - \sin \theta_1) \Re[\tilde{m}_y(\omega)]. \quad (\text{A.27})$$

Within linear response, the complex magnetization amplitude is related to the effective rf driving field h_{rf} (Spin-orbit field and Oersted field) by the acoustic-mode susceptibility,

$$\tilde{m}_y(\omega) = \chi_{\text{ac}}(\omega) h_{\text{rf}}. \quad (\text{A.28})$$

Substituting into [equation \(A.27\)](#) gives

$$V_{\text{SRE}} = \frac{\Delta R I_{\text{rf}}}{4} (\sin \theta_2 - \sin \theta_1) \Re[\chi_{\text{ac}}(\omega) h_{\text{rf}}]. \quad (\text{A.29})$$

Near the acoustic-mode resonance field H_{ac} , the susceptibility can be written as

$$\chi_{\text{ac}}(H) = \frac{C}{(H - H_{\text{ac}}) + i\Delta H_{\text{ac}}}, \quad (\text{A.30})$$

where C is a real constant and ΔH_{ac} is the linewidth. The real and imaginary parts

are

$$\Re[\chi_{ac}] = C \frac{H - H_{ac}}{(H - H_{ac})^2 + \Delta H_{ac}^2}, \quad (A.31)$$

$$\Im[\chi_{ac}] = C \frac{\Delta H_{ac}}{(H - H_{ac})^2 + \Delta H_{ac}^2}. \quad (A.32)$$

Here important to note that, additional field dependent AMR and GMR responses will be present in V_{SRE} response as a background and the presence of domains and their motion further complicates an accurate analysis of the SRE response for the eigenmodes. There will be also ANE present as background when the field is not collinear with the current ($\parallel \hat{x}$) direction.

A.2 GMR rectification for Optical mode

In the optical eigenmode, the two magnetizations precess out of phase, leading to an oscillation of the canting angle between the two magnetizations which leads to the rectification of the GMR response. The electrical resistance of a SAF due to GMR depends on the relative angle Φ between the two magnetization vectors:

$$R(\Phi) = R_p + \frac{\Delta R_{GMR}}{2} (1 - \cos \Phi). \quad (A.33)$$

For observing the optical mode solely, the magnetic field is applied perpendicular to the current direction ($\mathbf{j} \parallel \hat{x}$), in contrast to the acoustic mode. In the equilibrium state, the magnetizations lie in the x - y plane and are canted symmetrically with respect to the y axis:

$$\mathbf{m}_1 = (\sin \theta, \cos \theta), \quad (A.34)$$

$$\mathbf{m}_2 = (-\sin \theta, \cos \theta), \quad (A.35)$$

where θ is the canting angle of each magnetizations from the y -axis. The relative angle between the two magnetizations is therefore

$$\Phi_0 = \arccos(\mathbf{m}_1 \cdot \mathbf{m}_2) = 2\theta. \quad (A.36)$$

As the canting angle changes dynamically during the optical mode precession, we can define,

$$\theta(t) = \theta + \delta\theta(t), \quad (\text{A.37})$$

and hence a time-dependent relative angle

$$\Phi(t) = 2\theta + \delta\Phi(t), \quad \delta\Phi(t) = 2\delta\theta(t). \quad (\text{A.38})$$

For small oscillations $|\delta\Phi| \ll 1$, the resistance can be expanded to first order:

$$\delta R(t) = \left. \frac{dR}{d\Phi} \right|_{\Phi_0} \delta\Phi(t). \quad (\text{A.39})$$

From [equation \(A.33\)](#),

$$\frac{dR}{d\Phi} = \frac{\Delta R_{\text{GMR}}}{2} \sin \Phi. \quad (\text{A.40})$$

Evaluating at $\Phi_0 = 2\theta$ and using $\delta\Phi = 2\delta\theta$ yields

$$\delta R(t) = \Delta R_{\text{GMR}} \sin(2\theta) \delta\theta(t). \quad (\text{A.41})$$

The dynamic canting-angle response is described by the optical-mode susceptibility $\chi_{\text{opt}}(\omega)$:

$$\delta\theta(\omega) = \chi_{\text{opt}}(\omega) h_y^{\text{rf}}. \quad (\text{A.42})$$

In the time domain,

$$\delta\theta(t) = \Re[\chi_{\text{opt}}(\omega) h_y^{\text{rf}} e^{i\omega t}]. \quad (\text{A.43})$$

An rf current

$$I(t) = I_{\text{rf}} \cos(\omega t), \quad (\text{A.44})$$

flows through the device. The instantaneous voltage is

$$V(t) = I(t) \delta R(t). \quad (\text{A.45})$$

The rectified (dc) voltage is obtained by time averaging:

$$V_{\text{SRE}} = \langle V(t) \rangle. \quad (\text{A.46})$$

Substituting [equation \(A.41\)](#) and [equation \(A.43\)](#) and performing the time average gives

$$V_{\text{SRE}} = \frac{1}{2} I_{\text{rf}} \Delta R_{\text{GMR}} \sin(2\theta) h_y^{\text{rf}} \Re[\chi_{\text{opt}}(\omega)]. \quad (\text{A.47})$$

[Equation \(A.47\)](#) shows that the GMR rectification vanishes for the optical mode when $2\theta = 0$ or π , which corresponds to the parallel and antiparallel alignment of the two layers' magnetizations respectively.

Abbreviations

$\mathcal{PT}$	parity-time.
AF	antiferromagnet.
AFMR	antiferromagnetic resonance.
AHE	anomalous Hall effect.
AI	artificial intelligence.
AM	acoustic mode.
AMR	anisotropic magnetoresistance.
ANE	anomalous Nernst effect.
AP	anti-parallel.
BIA	bulk inversion asymmetry.
BLS	Brillouin light scattering.
CIP	current in plane.
CPP	current perpendicular to plane.
DL	<i>damping-like</i> .
DMI	Dzyaloshinskii–Moriya interaction.
DW	domain wall.
EA	easy axis.
FL	<i>field-like</i> .
FM	ferromagnet.
FMR	ferromagnetic resonance.

GHz	giga hertz.
GMR	giant magnetoresistance.
HDD	hard disk drive.
HM	heavy metal.
IEC	inter-layer exchange coupling.
iSGE	inverse spin galvanic effect.
ISHE	inverse spin Hall effect.
LHM	left-handed mode.
LLG	Landau-Lifshitz-Gilbert.
MFC	magnon frequency comb.
MFM	magnetic force microscopy.
MOKE	magneto-optical Kerr effect.
MRAM	magnetic random access memory.
MTJ	magnetic tunnel junction.
NM	non-magnetic.
OHE	ordinary Hall effect.
OM	optical mode.
ONE	ordinary Nernst effect.
p-MOKE	polar magneto-optical Kerr effect.
PHE	planar Hall effect.
PM	paramagnet.
PMA	perpendicular magnetic anisotropy.
PSD	power spectral density.
rf	radio frequency.
RHM	right-handed mode.
RKKY	Ruderman-Kittel-Kasuya-Yosida.
SAF	synthetic antiferromagnet.
SHA	spin Hall angle.

SHE	spin Hall effect.
SHNO	spin Hall nano-oscillator.
SIA	structural inversion asymmetry.
SNE	spin Nernst effect.
SOC	spin-orbit coupling.
SOT	spin-orbit torque.
SRE	spin rectification effect.
SSE	spin Seebeck effect.
STDP	spike time dependent plasticity.
STNO	spin-torque nano-oscillator.
STT	spin transfer torque.
SW	spin wave.
THE	topological Hall effect.
THz	tera hertz.
TMR	tunnel magnetoresistance.
tr-MOKE	time-resolved magneto-optical Kerr effect.

Scientific output

Publications

1. **P. K. Rout**, P. V. Prakash Madduri, S. K. Manna, and A. K. Nayak, *Field-induced topological Hall effect in the noncoplanar triangular antiferromagnetic geometry of Mn_3Sn* , *Physical Review B* **99**, 094430 (2019).¹⁶⁵
2. J. Godinho, **P. K. Rout**, R. Salikhov, O. Hellwig, Z. Šobáň, R. M. Otxoa, K. Olejník, T. Jungwirth, and J. Wunderlich, *Antiferromagnetic domain wall memory with neuromorphic functionality*, *npj Spintronics* **2**, 39 (2024).²⁰³
3. **P. K. Rout**, J. Godinho, F. Vilsmeier, R. Salikhov, J. A. Vélez, Z. Šobáň, D. Laroze, R. M. Otxoa, C. H. Back, O. Hellwig, and J. Wunderlich, *Spin-orbit torque driven antiferromagnetic oscillator*, *Under review*.
4. **P. K. Rout**, F. Vilsmeier, J. Godinho, R. Salikhov, Z. Šobáň, C. Riedel, O. Hellwig, C. H. Back, and J. Wunderlich, *Imaging of antiferromagnetic domain dynamics enabled by staggered spin-orbit fields*, *Under preparation*.

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