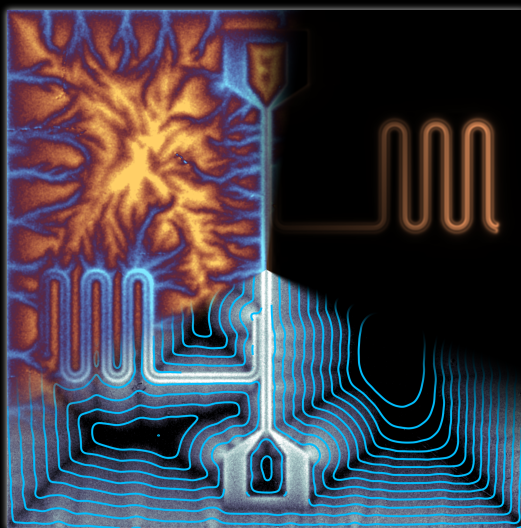

Mapping electric current and magnetic flux in structured superconductors



Nicolas LEJEUNE

ACADEMIC YEAR 2025-2026



FACULTY OF SCIENCES

QUANTUM MATERIALS RESEARCH UNIT

EXPERIMENTAL PHYSICS OF NANOSTRUCTURED
MATERIALS

Mapping electric current and magnetic flux in structured superconductors

AUTHOR: Nicolas LEJEUNE

SUPERVISOR: Prof. Dr. A.V. Silhanek

ACADEMIC YEAR 2025-2026

PhD thesis committee:*Chairman:*

Prof. Dr. N.D. Nguyen University of Liège

Secretary:

Prof. Dr. B. Vanderheyden University of Liège

Jury members:

Prof. Dr. C. Feuillet-Palma ESPCI Paris

Prof. Dr. D. Kölle University of Tübingen

Prof. Dr. J. Van de Vondel KU Leuven

Supervisor:

Prof. Dr. A.V. Silhanek University of Liège

Dissertation presented in partial fulfillment of the requirements for
the degree of *Doctor in Science*.

About the cover. The three panels illustrate the complementary imaging modalities used throughout this work. *Top left:* magnetic flux avalanches abruptly penetrating a superconducting resonator at low temperature, visualized by magneto-optical imaging. *Top right:* a simulated scanning laser microscopy image, mapping the current density distribution in a superconducting resonator. *Bottom:* a magneto-optical image recorded at higher temperature, where flux penetration is smooth and no avalanches occur; the superimposed lines are the electric current streamlines reconstructed from the measured field map by inversion of the Biot–Savart law. Each panel is an artistic rendering of one of the result chapters of this thesis.

Abstract

Superconducting planar resonators are essential elements of quantum computers and single-photon detectors. A growing number of technological applications require them to operate under an applied magnetic field. Unfortunately, under these conditions, this is where they present a fatal flaw. A perpendicular field of a fraction of a millitesla lets quantized magnetic flux enter the film, and the resonance frequency then responds to the field in an erratic, seemingly stochastic manner. Global microwave measurements, which provide an average response over the entire sample, have been unable to explain such response.

This thesis provides compelling evidence of the physical mechanism behind the observed stochastic phenomenon by probing locally the superconducting and magnetic responses, with micrometer resolution over millimeter scale field-of-view. Magneto-optical imaging (MOI) is utilized to map the magnetic flux threading the film. By inverting the Biot–Savart law we are then able to obtain the screening currents associated to the magnetic flux penetration. In addition to this technique, scanning laser microscopy (SLM) helps us to map a different parameter, the radio-frequency current of the resonant mode, which, in turns, reveals the locations where a resonator is most sensitive. The instrument used for the SLM technique was built from scratch in the course of this work. Next, each 2D map, whether it was obtained via MOI or SLM, is confronted with a single microwave measurement of the same device, so that a feature of the global response can be traced to a specific event occurring at a specific place.

On NbTiN resonators, the erratic response turns out to arise from

thermomagnetic flux avalanches. More precisely, four regimes of flux penetration are identified: complete magnetic screening, magnetic perforation of the ground plane, avalanche-dominated penetration and smooth flux removal. None of them is clearly distinguishable from the global microwave signal alone. By removing the metallic mirror of the magneto-optical indicator, we made it possible to record, for the first time, both measurements concurrently, and individual avalanches could then be tagged to specific jumps in resonance frequency. These measurements allowed us to conclude that the electrodynamic behavior is non-local: an avalanche occurring at a specific resonator, or in the feedline, shifts the resonance of every resonator on the chip. This also settles a question the literature had left open. Radio-frequency excitation in the linear Campbell regime only marginally enhances avalanche activity, precipitating avalanches rather than creating them.

Since high- T_c superconductors such as $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ (YBCO) are far less prone to develop flux avalanches, this material becomes an attractive alternative to conventional superconductors, and offers a further asset of its own when illuminated with visible light. In oxygen-deficient YBCO films, persistent photodoping is shown to be a slow process, stable and non-diffusive, in contrast to the fast and highly diffusive bolometric response exhibited at short times, providing a reversible and spatially selective way of modifying a resonator optically. Exploiting this contrast, a raster-scanned laser maps the local radio-frequency current density and reveals the standing waves at both the fundamental resonance and the second harmonic, identifying the regions that most strongly affect the resonance frequency. The mapping is then extended into the time domain: lock-in amplifiers are used to collect several hundred harmonics of a chopped photoresponse, and the Fourier series is inverted to reconstruct the time-dependence. We demonstrate that the resulting resolution of a few microseconds is set by the highest measurable harmonic and not by the acquisition speed. The resulting movies expose the delayed contribution of weakly current-carrying regions, governed by thermal diffusion.

Returning to conventional superconducting resonators, the tradi-

tional remedy has always been to keep the magnetic flux at bay with artificial pinning centers. Yet pinning sites are themselves known to favor avalanche triggering at low temperature. Laser-induced periodic surface structures (LIPSS) imprinted on Nb thin films offer a way out to this dilemma, because they can be instrumental to steer the penetrating flux instead of merely trapping it. Inverting the flux maps into screening-current maps shows that the submicron ripples produced by the LIPSS technique raise the critical current density along their ridges by a factor of three. The principal flux distributors, however, are not the ripples but the boundaries between orientation domains. Flux can accordingly be channeled, diverted, or stored in bundles of controllable size.

Taken together, these results trace a route from diagnosis to design. While SLM indicates where a resonator is most vulnerable to flux, laser-imprinted LIPSS offers the opportunity to guide the magnetic flux away from the sensitive zones.

Keywords: superconducting coplanar-waveguide resonators; magneto-optical imaging; scanning laser microscopy; thermomagnetic flux avalanches; RF-driven vortex dynamics; kinetic inductance; persistent photodoping; $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$; NbTiN; laser-induced periodic surface structures; magnetic flux guiding.

Résumé

Les résonateurs supraconducteurs planaires sont des éléments essentiels des ordinateurs quantiques et des détecteurs de photon unique. Un nombre croissant d'applications technologiques exige qu'ils fonctionnent sous champ magnétique appliqué. Or, c'est précisément dans ces conditions qu'ils présentent un défaut rédhibitoire. Un champ perpendiculaire d'une fraction de millitesla suffit à laisser pénétrer le flux magnétique quantifié dans le film, et la fréquence de résonance répond alors au champ de manière erratique, en apparence stochastique. Les mesures micro-ondes globales, qui fournissent une réponse moyennée sur l'ensemble de l'échantillon, n'ont pas été en mesure d'expliquer un tel comportement.

Cette thèse apporte des preuves convaincantes du mécanisme physique à l'origine du phénomène stochastique observé, en sondant localement les réponses supraconductrice et magnétique, avec une résolution micrométrique sur un champ de vue millimétrique. L'imagerie magnéto-optique (MOI) est utilisée pour cartographier le flux magnétique traversant le film. En inversant la loi de Biot–Savart, nous parvenons ensuite à obtenir les courants d'écrantage associés à cette pénétration de flux. À cette technique s'ajoute la microscopie laser à balayage (SLM), qui permet de cartographier une autre grandeur, le courant radiofréquence du mode résonant, laquelle révèle à son tour les régions où le résonateur est le plus sensible. L'instrument utilisé pour la technique SLM a été entièrement conçu et construit au cours de ce travail. Chaque carte 2D, qu'elle ait été obtenue par MOI ou par SLM, est ensuite confrontée à une unique mesure micro-onde du

même dispositif, de sorte qu’une signature de la réponse globale puisse être rattachée à un événement précis survenant en un endroit précis.

Sur des résonateurs en NbTiN, la réponse erratique s’avère provenir d’avalanches thermomagnétiques de flux. Plus précisément, quatre régimes de pénétration du flux sont identifiés : écrantage magnétique complet, perforation magnétique du plan de masse, pénétration dominée par les avalanches et évacuation régulière du flux. Aucun d’eux n’est clairement discernable à partir du seul signal micro-onde global. En retirant le miroir métallique de l’indicateur magnéto-optique, nous avons rendu possible l’enregistrement simultané des deux mesures, pour la première fois, et chaque avalanche a pu dès lors être associée à un saut spécifique de la fréquence de résonance. Ces mesures nous ont permis de conclure que le comportement électrodynamique est non local : une avalanche survenant dans un résonateur donné, ou dans la ligne d’excitation, décale la résonance de tous les résonateurs de la puce. Cela tranche également une question restée ouverte dans la littérature : l’excitation radiofréquence dans le régime de Campbell linéaire n’accroît que marginalement l’activité d’avalanches, celles-ci étant précipitées plutôt qu’engendrées.

Les supraconducteurs à haute température critique tels que le $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ (YBCO) étant bien moins enclins à développer des avalanches de flux, ce matériau devient une alternative attrayante aux supraconducteurs conventionnels, et offre un atout supplémentaire qui lui est propre lorsqu’il est illuminé en lumière visible. Dans les films de YBCO désoxygénés, nous montrons que le photodopage persistant est un processus lent, stable et non diffusif, par contraste avec la réponse bolométrique rapide et fortement diffusive observée aux temps courts, ce qui offre un moyen réversible et spatialement sélectif de modifier optiquement un résonateur. Exploitant ce contraste, un laser balayé point par point cartographie la densité locale de courant radiofréquence et révèle les ondes stationnaires aussi bien à la résonance fondamentale qu’à la seconde harmonique, identifiant les régions qui influencent le plus fortement la fréquence de résonance. La cartographie est ensuite étendue au domaine temporel : des am-

plificateurs à détection synchrone sont utilisés pour collecter plusieurs centaines d'harmoniques d'une photoréponse modulée par hacheur, et la série de Fourier est inversée afin de reconstruire la dépendance temporelle. Nous démontrons que la résolution ainsi obtenue, de quelques microsecondes, est fixée par la plus haute harmonique mesurable et non par la vitesse d'acquisition. Les séquences temporelles qui en résultent mettent en évidence la contribution retardée, gouvernée par la diffusion thermique, des régions faiblement porteuses de courant.

Revenant aux résonateurs supraconducteurs conventionnels, le remède traditionnel a toujours consisté à tenir le flux magnétique à distance au moyen de centres de piégeage artificiels. Or les sites de piégeage sont eux-mêmes connus pour favoriser le déclenchement des avalanches à basse température. Les structures périodiques de surface induites par laser (LIPSS) imprimées sur des films minces de Nb offrent une issue à ce dilemme, car elles peuvent servir à orienter le flux pénétrant au lieu de se contenter de le piéger. L'inversion des cartes de flux en cartes de courants d'écrantage montre que les ondulations submicrométriques produites par la technique LIPSS augmentent d'un facteur trois la densité de courant critique le long de leurs crêtes. Les principaux distributeurs de flux ne sont toutefois pas les ondulations elles-mêmes, mais les frontières entre domaines d'orientation. Le flux peut dès lors être canalisé, dévié ou stocké sous forme de paquets de taille contrôlable.

Pris ensemble, ces résultats dessinent un cheminement du diagnostic vers la conception. Tandis que la SLM indique où un résonateur est le plus vulnérable au flux, les LIPSS imprimées par laser offrent la possibilité de guider le flux magnétique loin des zones sensibles.

Mots-clés : résonateurs supraconducteurs à guide d'onde coplanaire ; imagerie magnéto-optique ; microscopie laser à balayage ; avalanches thermomagnétiques de flux ; dynamique des vortex sous excitation RF ; inductance cinétique ; photodopage persistant ; $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$; NbTiN ; structures périodiques de surface induites par laser ; guidage du flux magnétique.

Acknowledgments

Writing a thesis offers you a one-of-a-kind moment: the moment to sit back and look at your life to really appreciate how well surrounded you are. These few (I hope not too long) pages have the modest goal of acknowledging all the people who have crossed my path and made my journey through these past four years a bit more enjoyable/easier to handle.

There's no doubt about who should be acknowledged first: thank you Alejandro! From the day we started (really) working together¹, you have always been a driving force. I still don't understand how a simple meeting with you could make me hop from the burnout state to the over-motivated one (one could naively think this transition is forbidden). Thank you for all the times you trusted me even when I was not trusting myself, for the technical freedom you gave me, for the daily guidance you share, and for always being available for a quick (and sometimes not so quick) chat whenever I needed it. You have really been a mentor to me, and I hope this thesis makes you proud of the work we accomplished together!

Next, I would like to acknowledge those who (I hope) will deliver me from the status of PhD student, namely my PhD Jury and more specifically Prof. Dieter Kölle, Prof. Cheryl Feuillet-Palma, Prof. Joris Van de Vondel, Prof. Benoit Vanderheyden and Prof. Ngoc Duy Nguyen. I hope that you will enjoy reading this thesis, and that I was

¹It indeed took a few months before the two socially awkward people we are understood each other.

somehow able to put on paper the excitement I had when conducting all those experiments.

Two members of the Jury deserve a special mention. If Alejandro is my academic father, then Duy Nguyen, who heads the other half of the lab, must be my academic uncle: he trusted me with the lab equipment we share at the Nanofab platform, and again when it came to teaching the Electronics course. Joris Van de Vondel is one of the protagonists of the longstanding collaboration between Liège and Leuven in which he welcomed me from the very beginning of my PhD. Every single visit to Leuven was a pleasure, both scientifically and socially!

Talking about collaboration, this thesis owes a great deal to those who contributed to the work it reports. Prof. Antonio Badía-Majós offered me wonderful exchanges during the investigation of the LIPSS sample, and welcomed me for a research stay. I enjoyed meeting him and his team, and I feel like it really pushed the collaboration forward! The team of Dr. Anna Palau was equally essential to what was achieved here. And even though I have already talked about Leuven, Lukas Nulens deserves a line of his own: a great and very nice scientist to work with.

I am greatly indebted to the FNRS which, through the FRIA grant they awarded me, gave me the priceless opportunity to pursue research in a prolific environment. I am aware of how lucky I was to attend so many conferences and research stays, and to meet so many scientists from all over the world.

Along the way, I've had quite a few colleagues who were either great collaborators, lunch partners or simply friends. Emile stands out among them, for talking me out of a career in theoretical materials (a fate I was probably unfit for anyway) and for passing on to me his passion for cowboy-style science. I don't forget the colleagues who have already left to follow their own path: Thomas R., Suraj, Simon C., Naama, Stefan; nor those who still share it with me for a short while longer: François, Thomas P., Abhishek, Daniel, Cyril, Julia, Caio, and finally my last Master's student Simon VdH, who

I'm convinced will produce excellent work with the SLM setup². Dear colleagues, never forget to keep a few numbers available in your pocket of numbers; it will serve you someday. Please also play loud some Goldman in the Cryo lab from time to time, so that the cryostat remembers the good times we've had together. One colleague³ is missing from the list, because he's the one with whom there's the largest time overlap. Amaury Baret and I started the Bachelor in Physics together, and we (nearly) finish at the PhD stage together. It was a real pleasure to share (too many) breaks discussing science, life (and gossip). I hope someday in the future we'll meet at conferences, and I'm looking forward to the day we'll finally be able to publish together. Thanks Amaury!

The rest of the acknowledgments will switch to the appropriate language (Abhishek, this is one of the last occasions you have to learn French; same to you for Polish, Julia).

L'Université de Liège se doit de fonctionner comme une horloge au quotidien, et en horlogerie, chaque engrenage est essentiel. Je saisis donc l'opportunité de prendre quelques lignes pour remercier ceux que je côtoyais presque tous les jours. Merci Michel pour ta sympathie et ta disponibilité à toute heure du jour et de la nuit, merci Florence pour ta gentillesse lors des centaines de commandes pas toujours claires que nous avons passées ensemble, et merci Alexis pour ton accueil toujours souriant !

Ensuite, je dois reconnaître que j'ai particulièrement de la chance d'être entouré de parents qui m'ont tant apporté. Merci maman pour toutes les fois où tu as fait de tes enfants ta priorité ; j'espère que tu es fière de ce que nous sommes devenus. Merci papa pour les encouragements (parfois maladroits) mais toujours sincères ! Merci aussi à Marie et Laurent, qui ont fait le même chemin avant moi et m'ont convaincu que c'était possible. Même sans grands discours, vous savez toute la gratitude et l'affection que je vous porte. Granny,

²I hope I was able to teach you that the strategy "Act first, think next" is efficient.

³I hope only one.

elle, n'a pas pu rester jusqu'à la fin du chemin que je parcourais. Elle qui a toujours rêvé d'avoir un petit-enfant médecin, j'espère que, de son étoile, elle saura faire l'amalgame et qu'elle est fière d'avoir un nouveau docteur dans la famille. Na koniec chciałbym serdecznie podziękować mojej babci, która zawsze wierzyła we mnie najbardziej. Bardzo się cieszę, że jest przy mnie już tak długo. Kocham cię, babciu.

Je tiens également à exprimer ma gratitude envers ma professeure de physique du secondaire, Martine Vertez, qui pendant trois ans a su éveiller ma curiosité scientifique et me donner l'envie initiale de me lancer dans la physique.

Les derniers, mais pas les moindres, sont ceux qui m'ont permis tant de fois de décrocher de la thèse. Mes amis et amies sont ma famille de cœur, et j'espère qu'ils et elles se rendent compte du pilier essentiel qu'ils et elles représentent. Je prends le risque de les citer, bien que ceux que je pourrais oublier ne soient oubliés que sur cette page. Merci à Danaé et Diane, Adrien, Colin (café ?), Simon, Maxime, Fred, Robin, Bruno, Pol, Julie, Jaco, Hélène, Eline, France, Alix, et tous les autres qui se reconnaîtront !

Je ne me souviens que d'un mur immense

Mais nous étions ensemble

Ensemble, nous l'avons franchi

— Jean-Jacques Goldman, *Ensemble*

Statement about the use of generative AI

Nowadays, a statement about the use of AI deserves, in my opinion, more than a single sentence copied and pasted from one thesis to the next. The question is recent and still unsettled, and I would rather give it a proper explanation than a formulaic disclaimer; that is what follows.

The thesis I am now defending is based on work carried out between the middle of August 2022 and today, at the end of July 2026. This context is essential to understanding how AI was used in the writing process. When this work began, generative AI was not yet a widely discussed subject, as the pioneer of the field, OpenAI, had not yet released ChatGPT. Everything changed on 30 November 2022, when ChatGPT was presented to the world and people could begin to experience what it felt like to converse with an agent that answered questions with some form of memory, something fundamentally different from the way Google and other assistants handled tasks. Personally, I remember being amazed at how ChatGPT would attempt any task and could deal with subjects Google was incapable of handling. To be fair, at that time few people were using it for genuinely useful purposes: we were essentially asking it to make up stories with outlandish plots, or to tell us how much a *baudet du*

Poitou⁴'s egg might weigh (according to early ChatGPT, it was remarkably light, between 150 and 200 g). Through these interactions, we were discovering that the AI was designed to answer the user no matter the question, which occasionally led to what was later termed *hallucinations*: answers entirely fabricated to satisfy the user's request. As a result, generative AI was seen more as an entertaining curiosity than as a useful tool.

Then the years passed, and more and more of today's AI giants began competing to release the best generative model. Gradually, hallucinations became less frequent, and people started using these tools for simple everyday tasks such as drafting email replies, preparing shopping lists, or inventing recipes from whatever was left in the fridge. Their use nevertheless remained frowned upon in science, since academic tasks would most often produce answers containing made-up citations and DOIs, claims unsupported by the literature, and so on.

I remember it was at the end of the third year of my PhD that the agent developed by Google, Gemini, with its thinking feature, began to give scientifically grounded and physically argued responses. Even so, because I had witnessed the early stages of AI, I always triple-checked whatever came out of the machine, validating the method it used to solve typical problems. With that critical mindset, and like a growing number of scientists around the world, I began using Gemini as a kind of coffee-break discussion partner. It became what we call in our lab a *scientific punching ball*: you overwhelm it with ideas, and while trying to convince it that you grasp the underlying physics, you start to see the weak points of your own theory and identify clearly what still needs to be investigated, or where your reasoning needs to be revised.

Today, it would be naive to think that generative AI is still at the level of its early versions. Over three or four years, I have personally watched its capabilities grow enormously, and it is now an extremely powerful tool, provided one always remembers that, as has always

⁴A French donkey breed, 1.5 m tall and weighing roughly 400 kg, that obviously does not lay eggs.

been the case in science, critical thinking is essential when confronting new claims. There remains, and I can still clearly sense it within the scientific community, a kind of taboo around the use of AI. Few dare to state openly that “AI was used to improve the quality of the text” or that “AI was employed to find new ideas or new ways of tackling a problem”. In my view, this reluctance is a mixture of two things. Part of it is a perfectly legitimate distrust: given the fabricated citations and unsupported claims that AI has produced, and can still produce, caution is a sound scientific reflex rather than mere conservatism. But part of it is also the defensive reaction humans tend to display whenever a new technology arrives, a certain technophobia rooted in the fear of being replaced, which leads people to shun the tool rather than confront it (the same happened, for instance, with automatic spell-checking, which people were long reluctant to admit using).

This thesis was written between April 2026 and July 2026, at a time when AI is unambiguously a powerful tool when used properly. AI was therefore, unsurprisingly, used to help me write it, but only very cautiously and at different levels, which are detailed below.

The theoretical Chapters 2 and 3 are conceived as a condensation of the essential physical concepts needed to address the rest of the thesis. They are the chapters in which generative AI was used most, but this reflects the nature of the task rather than any greater contribution of the tool to their substance. Being purely expository, they report no original result of my own: the whole writing task there consists in articulating well-established physics into a clear and continuous narrative, that is, precisely the connective prose I had chosen to delegate. Its role was therefore confined entirely to the phrasing of sentences whose content I had already fixed. I first decided the complete structure of each chapter, every physical concept to be developed and how to do it, the logical links between them, and the order in which they had to appear, and only then did I delegate the writing to generative AI⁵, with a prompt aimed solely at turning that fixed structure into connected sentences, and explicitly forbidding the

⁵The tool used for these two chapters was Claude (Anthropic).

introduction of any physical concept, claim, or idea I had not planned.

I was, throughout, acutely aware of the characteristic failure mode of such tools: even when asked only to rephrase, a language model will readily insert plausible-sounding but unrequested physics, spurious justifications, or subtly wrong statements presented with full confidence. My repeated proofreading of these chapters was therefore not a cosmetic pass but an active filtering against exactly this risk, sentence by sentence, discarding or correcting anything that departed from what I had decided to convey. All figures are entirely my own creation. As a result, I fully endorse and take responsibility for the content of these chapters: the physics they present is mine alone, and each formulation was checked against my intent before being kept. I therefore stand behind them as I would behind any text I had written word for word myself.

Chapter 4 was written entirely by me, with AI used only as an English-reformulation tool. This is the use most people are nowadays beginning to admit to, and it does not fundamentally differ from the use described above: the core of the text comes from the author, and AI merely shapes a solid, grammatically correct body around it.

Finally, Chapters 5, 6 and 7 are adapted from existing publications, as stated in their respective introductions. Their text therefore comes directly from work that was written and peer-reviewed before this thesis, and no generative AI was used in producing them.

To summarize this long but essential attempt to put generative AI in perspective: it should be regarded as nothing more than a tool, just as spell-checkers, calculators, spreadsheets, high-level programming languages, and Wikipedia are tools that help their users become more productive, with the added responsibility, for the user, of vouching for the accuracy of what is produced. I am deeply convinced that in the years to come this tool will revolutionize the way science is done, but also that it will become widely accepted that it is always the scientist who decides what is worth keeping and what amounts to a hallucination.

Contents

Abstract	v
Résumé	ix
Acknowledgments	xiii
1 Introduction	1
2 Theory	9
2.1 Superconducting microwave resonators	9
2.1.1 From transmission line to resonant cavity . . .	11
2.1.2 Distributed parameters: geometry and substrate effects	14
2.1.3 Microwave characterization	16
2.1.4 Kinetic inductance: a uniquely superconducting contribution	19
3 MOI: Principles and Practices	25
3.1 Physical principle	25
3.2 Experimental setup	28
3.3 Post-processing	30
3.3.1 Calibration: intensity to magnetic field	30
3.3.2 MOI inversion: from flux density to current density	31

4	Scanning Laser Microscopy	35
4.1	Experimental setup	36
4.2	Detection modes	37
4.2.1	Optical mode	39
4.2.2	Electrical modes	41
4.3	Periodic time reconstruction	51
5	Magnetic Flux in Low-T_c Resonators	55
5.1	Sample design	58
5.2	Experimental methods	59
5.2.1	Sample fabrication	59
5.2.2	RF measurements	61
5.2.3	Magneto-optical imaging	62
5.2.4	SQUID magnetometry	62
5.3	Temperature-dependent resonator characterization . .	62
5.4	Magnetic field response	63
5.4.1	Different regimes of flux penetration	63
5.5	Temperature dependence of avalanches	67
5.5.1	Strategies for magnetic field resilience	69
5.6	Simultaneous RF and MOI measurements	70
5.6.1	Characterization of Device #2	71
5.6.2	Invasiveness of the magneto-optical indicator .	73
5.6.3	Effect of RF excitation on the avalanche activity	76
5.6.4	Effect of avalanche activity on the RF transmis- sion	78
5.6.5	Electromagnetic simulations	81
5.6.6	Dynamic regime analysis	85
5.7	Conclusions	86
6	SLM of High-T_c YBCO Devices	89
6.1	Photoexcitation of YBCO resonators	91
6.1.1	Introduction	91
6.1.2	Sample fabrication and resonator design . . .	93
6.1.3	Results and discussion	96

6.1.4	Intermediate conclusion	105
6.1.5	Annex	105
6.2	Time-resolved SLM	113
6.2.1	Introduction	113
6.2.2	Case study: superconducting resonator	114
6.2.3	Conclusion	121
7	Steering Magnetic Flux with LIPSS	123
7.1	Introduction	125
7.2	Materials and Methods	129
7.2.1	Fabrication	129
7.2.2	Microstructural characterization	132
7.2.3	Magneto-optical imaging	133
7.2.4	Magnetic characterization	133
7.3	Critical-state modeling	134
7.4	Results	136
7.4.1	Tailoring Nb thin films: laser-processed meta- surfaces	136
7.4.2	Routing magnetic flux penetration	137
7.5	Discussion	147
7.6	Conclusions	152
8	Conclusions and outlooks	157
8.1	Summary of the main results	157
8.2	Laser-machined resonators	161
8.3	Experimental limitations	164
8.4	Outlook	166
	Bibliography	169
	Curriculum Vitae	201

Chapter 1

Introduction

A superconducting resonator is, in essence, a piece of superconducting thin film patterned into a circuit that resonates at a well-defined frequency, in the same way that a guitar string rings at a well-defined pitch that depends on its length. Because a superconductor carries current without resistance, such a circuit rings for an extraordinarily long time: the quality factor Q , which quantifies the number of oscillations before the energy has appreciably decayed, can reach 10^7 , several orders of magnitude beyond what any normal metal allows [1]. That sharpness is what makes these devices useful. A resonance so narrow shifts measurably in response to a minute perturbation, and the same property that makes them exquisite sensors also makes them the building blocks of superconducting qubits. They are today essential elements of quantum computing architectures [2–4], of single-photon detectors [5], and of quantum sensing and information processing [6].

Their weakness is equally well known, and it is magnetic. A type-II superconductor expels a magnetic field only up to a threshold field beyond which the field enters the material not continuously but in discrete threads, each carrying one quantum of magnetic flux. These flux quanta, or *vortices*, are the source of the trouble [7]. Driven back and forth by the radio-frequency (RF) current circulating in the device, they dissipate energy, which lowers the quality factor, and they

alter the way the superconductor stores magnetic energy, which shifts the resonance frequency [8–11]. In a thin film, a perpendicular field of a fraction of a millitesla is already enough for vortices to enter [12].

One might simply consider shielding the device, as is often done. However, the presence of a magnetic field (in-plane or out-of-plane) is a necessary condition for a growing number of applications. For instance, spin ensembles in solid-state systems, which require 15–200 mT [13–15], phase-slip devices [16–18], and trapped electrons, which require fields in the tesla range [19, 20]. Even when the field is nominally applied in the plane of the film, where these circuits are most resilient [21], a slight misalignment suffices to produce a substantial perpendicular component. Making superconducting resonators resilient to magnetic field is therefore a technological problem in its own right. Two decades of work have been devoted to it along two broad lines: hindering vortex motion with artificial pinning sites [9, 11, 22], or reducing the number of vortices in the critical regions by narrowing or fractalizing the ground plane [23–25].

A mysterious response, and a blind spot. None of these strategies, however, take into account a phenomenon that global measurements have never been able to explain. When the resonance frequency of a planar resonator is recorded under an increasing perpendicular magnetic field, it does not vary smoothly. Instead, it jumps, erratically and irreproducibly, in a manner that looks like noise and has often been treated as such [26–28].

The reason why this response has resisted explanation is consequence of a blind spot in the characterization techniques. The most common ones, such as electrical transport, magnetometry, resonance frequency, each return a single number, averaged over the entire sample. Yet the underlying phenomena are intrinsically spatially local. Flux does not seep in uniformly: it enters through a handful of preferred nucleation points. Screening currents do not spread evenly: they crowd at the edges. A global measurement can establish *whether* a device misbehaves, but can not tell *why*, and it therefore cannot

help to identify which part of the device should be re-engineered.

This thesis removes that blind spot, and in doing so attributes the erratic response to a specific mechanism: *thermomagnetic flux avalanches*. Below roughly half the critical temperature T_c , the heat released by moving vortices locally deteriorates the screening ability, which lets more flux in, which in turn releases more heat. This runaway floods a large fraction of the film with flux in nanoseconds [29], and each such event displaces the resonance frequency abruptly. What looks like noise is in fact a sequence of discrete, localizable, physically intelligible events.

Identifying the mechanism immediately raises three questions, which are addressed throughout this thesis.

- ***Does the radio-frequency excitation inevitably provoke the avalanches?*** The literature is contradictory: some experiments report RF-stimulated avalanches [30–32], others report suppression near the depinning frequency [33], and simulations predict a reduced threshold field under AC excitation [34]. The disagreement is not accidental. No experiment had witnessed the flux and the resonance at the same time, so avalanche activity had to be inferred from the very microwave trace whose interpretation was in question.
- ***Which portion of a resonator determines its resonance frequency?*** One would rationally assume that the response is modified most strongly in the regions where the standing-wave current is the highest, and that these regions are set by the boundary conditions of the resonator. However, this has seldom been experimentally measured. This question is directly linked to the first one, as it identifies the critical location where a vortex avalanche would exert the most detrimental impact.
- ***Is keeping flux out the only option?*** Every strategy described above relies on exclusion or immobilization of flux quanta, and the most common of them actually promote magnetic flux

avalanches. Indeed, pinning sites reduce flux-shaking losses at high temperature. However, at low temperature they are known to *favor* avalanche triggering, and so to degrade the very device they were meant to protect [35, 36]. Whether penetrating flux can instead be actively *directed*, sent along chosen paths and away from the regions that matter, is an open question we will attempt to tackle.

Local probes of a global response. All three questions above concern spatially localized events, something that a global measurement cannot discern. The strategy adopted throughout this thesis is accordingly the same in every chapter: map, with micrometer resolution over millimeter-wide fields of view, the local parameter that governs the phenomenon at hand, and confront that map with the averaged observable that a conventional measurement would return. The relevant local parameter changes from chapter to chapter, and so does the microscope used to measure it.

Three parameters are relevant to the investigated systems. The first is the *magnetic flux density* threading the film, which reports on field penetration, vortices and avalanches. The second is the *screening current*, the direct current circulating in the film in response to the applied magnetic field, first in the magnetic screening state and then in the critical state. This current determines where flux enters and how far it advances. Currents and magnetic fields are rigidly tied since in a thin-film geometry, the perpendicular flux density at any point is fixed by the screening-current distribution across the whole sample, and conversely [37]. The third is the *radio-frequency current* of the resonant mode, and it is of an altogether different nature. Set by the circuit geometry and by the standing-wave pattern of the mode, it oscillates at gigahertz frequencies and is present even in zero applied magnetic field. This RF current dictates where a resonator is sensitive, and hence where a vortex or an avalanche is most harmful.

Two techniques give access to these quantities. *Magneto-optical imaging* (MOI) exploits the rotation of the polarization of light in

a magnetic field. A thin magneto-optical crystal laid on the sample converts the local out-of-plane flux density into an optical intensity, giving a direct picture of where the flux is [30, 38]. Inverting the Biot–Savart law then allows one to recover the screening currents that produced that picture. *Scanning laser microscopy* (SLM) addresses the RF current, and does so indirectly. A focused laser beam is scanned across the device, warming one micrometer-sized spot at a time; the perturbation caused in the global microwave response reveals how strongly that spot participates in the resonant mode [39, 40].

Crucially, neither technique is used here as a stand-alone imaging experiment. Each is combined with a microwave measurement of the device it images, so that a feature appearing in the global response can be assigned to an identified event, or an identified region, on the map. That coupling is what turns an image into an explanation, and it is the methodological thread of the whole manuscript.

What the maps reveal. Applied to a low- T_c NbTiN resonator, this treatment dissolves the noisy field response into four distinct regimes: screening, a magnetic perforation of the ground plane, avalanche-dominated penetration, and a smooth flux removal that ends in *anti-avalanches*, sudden entries of flux of polarity opposite to the applied field. None of these regimes is clearly discernible from the microwave signal alone, and the correspondence extends down to details that would otherwise be dismissed as noise. A small kink in the resonance frequency, for instance, turns out to mark the precise moment at which avalanches begin in the ground plane of the resonator meander rather than in the borders of the sample. The agreement is close enough that the argument can be run backwards: the resonator becomes a sensitive and non-invasive detector of the flux dynamics in its own film.

Driving the two measurements simultaneously rather than in sequence requires adapting the optical cryostat to perform simultaneous RF characterization measurements. It sharpens the probing down to an event-by-event correspondence, in which individual avalanches are

tagged to individual frequency jumps, and electromagnetic simulations relate the size of a jump to where the avalanche landed. This is what makes the first question answerable at all: with flux and resonance recorded on the same time axis, one no longer has to infer avalanche activity from the trace under scrutiny.

As high- T_c materials are far less prone to avalanches, $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ (YBCO) presents itself as an attractive alternative platform, and it turns out to offer a distinctive asset of its own. Oxygen-deficient YBCO can be *photodoped*: prolonged illumination improves its superconducting properties persistently and reversibly, an effect that is slow and stationary whereas ordinary laser heating is fast and diffusive. That contrast is exploited here in two ways. It provides a means of modifying a resonator optically, in contrast to the enhancement obtained by heavy-ion irradiation, which is permanent and destroys the device for any subsequent comparison [41]. In addition, it allows the second question above to be answered, by mapping the RF current of the resonant mode and thereby identifying which regions of a device govern its resonance frequency. No instrument capable of probing the latter existed in Liège at the start of this work, and this thesis accordingly builds one from scratch with that precise objective in mind. That construction spans the structural design and the choice of optical components, the assembly, the control software, and the calibration of the spatial and thermal response.

The third question is taken up on a different platform. Laser-induced periodic surface structures (LIPSS), submicron ripples imprinted on a niobium film by femtosecond laser pulses, do not merely trap flux where it is: they *steer* it, guiding it along or away from chosen regions. Here the flux maps are inverted to recover the screening currents themselves, which reveals that the ripples make the film electrically anisotropic and explains why flux follows the paths it does. Exclusion and immobilization are therefore not the only options, and the pinning dilemma can be sidestepped. Flux that cannot be kept out may instead be routed away from where it does the most damage, and identifying such places is exactly what the sensitivity maps above

provide.

A single connecting line threads through these chapters. Spatial mapping first *reveals* a mechanism hidden inside an averaged signal, then *localizes* the regions in which that mechanism is most harmful, and finally opens a route to *control* it. The recurring question, posed anew each time, is how the spatial configuration of the relevant local quantity explains, and ultimately governs, the measured global response.

Structure of the thesis. The manuscript is organized into a theoretical foundation, two methodological chapters, three experimental chapters, and a conclusion.

Chapter 2 lays out the theory needed to understand the rest of the thesis: the physics of superconducting resonators and the notions of kinetic inductance and quality factor, through which the local physics ultimately manifests itself in a measured microwave response.

Chapters 3 and 4 introduce the two mapping techniques and the developments and refinements made to them during this thesis. Chapter 3 describes magneto-optical imaging, covering the Faraday-rotation principle, the indicators, and the Biot–Savart inversion that converts a flux map into a screening-current map. Chapter 4 describes the scanning laser microscope. This chapter also presents the multiharmonic scheme developed to operate the microscope in the time domain. Both chapters are written as practical guides to the tools rather than as exhaustive reviews of the techniques, although in enough detail for either setup to be reproduced.

The three experimental chapters then apply these tools to address the questions above. Chapter 5 combines MOI with microwave transmission on low- T_c NbTiN resonators to witness flux entering the device while simultaneously tracking the resonance. Chapter 6 turns to high- T_c YBCO resonators and uses SLM in two ways: first to exploit persistent photodoping as a means of locally and reversibly modifying a resonator while mapping its RF current density, then to resolve that response in time and expose its dynamics and non-

locality. Chapter 7 takes the step from diagnosis to design, showing by MOI that laser-induced periodic surface structures imprinted on a superconducting film can actively *steer* the penetrating magnetic flux, guiding it along specific paths and away from chosen regions.

Chapter 8 draws these threads together. The steering demonstrated with LIPSS and the sensitivity maps obtained by SLM point back to the flux-avalanche problem of Chapter 5: SLM can indicate *where* a resonator is most vulnerable to flux, and LIPSS suggests *how* flux might be diverted from precisely those regions, sketching a route toward magnetic-field-resilient resonators. A first laser-machined niobium resonator, fabricated and read out with the very tools presented here, is reported as a preliminary step in that direction.

A word on the boxes. Chapters 2, 3 and 4 make frequent use of shaded *theory boxes*. Each box is a self-contained treatment of one supporting result, a derivation, a model, or a calibration, such as the telegrapher’s equations, the origin of kinetic inductance from the London equations, or the Biot–Savart inversion for thin-film current imaging. They are deliberately set apart from the main text so that the narrative can be followed without interruption: a reader already familiar with a given result may skip its box, while a reader wanting more detail will find it there. Later chapters refer back to these boxes rather than repeating their contents.

Chapter 2

Theory

2.1 Superconducting microwave resonators

Superconducting microwave resonators are planar microwave cavities patterned from superconducting thin films. Their defining advantage over normal-metal cavities is zero direct-current electrical resistance, which allows electromagnetic energy to be stored with negligible ohmic dissipation. The resulting internal quality factors Q_i are extraordinary: values above 10^6 are routinely achieved in optimized aluminum or titanium nitride devices [42, 43], with the best current planar resonators reaching $Q_i \approx 10^7$ [1]. To appreciate what this means, a quality factor of 10^6 at a resonance frequency of $f_0 = 5$ GHz corresponds to a photon lifetime $\tau = Q_i/(2\pi f_0) \approx 32 \mu\text{s}$, and to a resonance linewidth (full width at half maximum) $\Delta f = f_0/Q_i = 5$ kHz. Frequency shifts of this order, superimposed on a multi-gigahertz carrier, are straightforwardly resolved with a standard vector network analyzer (VNA), making superconducting resonators remarkably sensitive probes of any perturbation that modifies their electromagnetic environment.

This sensitivity has driven two broad fields of application. In quantum information science, superconducting resonators are used as quantum buses that couple distant superconducting qubits and as readout elements for their quantum states, operating at the single-

photon level [2, 44, 45]. In astrophysics and photon sensing, kinetic inductance detectors (KIDs) [5, 46] exploit the fact that absorption of a single high-energy photon successively breaks thousands of Cooper pairs, causing a measurable shift of the resonance frequency of a microwave resonator. In condensed matter physics, superconducting resonators have been used to detect individual vortex events, thermomagnetic instabilities, and changes in the local superfluid density caused by light illumination [28, 47, 48], all of which are at the core of this thesis.

The specific geometry employed throughout this work is the coplanar waveguide (CPW) resonator [49, 50]: a central metallic strip of width w , separated from two infinite coplanar ground planes by gaps of width g , all patterned in the same plane on top of a dielectric substrate (Fig. 2.1). The CPW supports a quasi-transverse electromagnetic (quasi-TEM) guided mode whose properties are entirely determined by the conductor geometry and the dielectric environment.

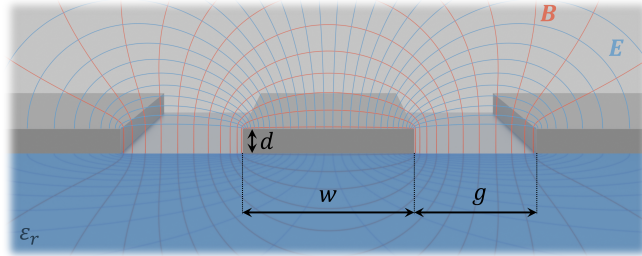


Figure 2.1: Cross-sectional schematic of a coplanar waveguide. A superconducting film (thickness d) is patterned on a substrate of dielectric constant ϵ_r and thickness $h \gg w$. The central conductor of width w is flanked by two gaps of width g and two semi-infinite ground planes. The RF electromagnetic field is concentrated in the gap regions, distributed between air (above) and substrate (below).

2.1.1 From transmission line to resonant cavity

A CPW is a distributed-parameter transmission line: rather than lumped inductance and capacitance, it possesses an inductance L' and a capacitance C' *per unit length*, which jointly govern the propagation of voltage and current waves along the structure [51, 52]. Their values are set by the conductor geometry and the surrounding dielectric media, as discussed in Section 2.1.2. Box 2.1 summarizes the wave propagation results that follow from them.

Box 2.1 – Telegrapher's equations and wave propagation

A transmission line with distributed inductance L' and capacitance C' per unit length obeys the *telegrapher's equations* [51, 52]:

$$\frac{\partial V}{\partial x} = -L' \frac{\partial I}{\partial t}, \quad \frac{\partial I}{\partial x} = -C' \frac{\partial V}{\partial t}. \quad (2.1)$$

These combine into a wave equation $\partial^2 V / \partial x^2 = L' C' \partial^2 V / \partial t^2$, whose general solution is a superposition of forward- and backward-propagating voltage waves:

$$V(x, t) = V_+ e^{i(\omega t - kx)} + V_- e^{i(\omega t + kx)}, \quad (2.2)$$

with wavenumber $k = \omega / v_{\text{ph}}$, $i^2 = -1$ and **phase velocity**

$$v_{\text{ph}} = \frac{1}{\sqrt{L' C'}} = \frac{c}{\sqrt{\varepsilon_{\text{eff}}}}, \quad (2.3)$$

where c is the speed of light in vacuum and ε_{eff} is the effective permittivity of the guiding structure (see Box 2.2). The corresponding current wave satisfies $I_{\pm} = \pm V_{\pm} / Z_0$, where the **characteristic impedance** is

$$Z_0 = \sqrt{\frac{L'}{C'}}. \quad (2.4)$$

In practice, CPW resonators are designed for $Z_0 = 50 \, \Omega$ to match standard microwave instrumentation. Equations (2.1)–(2.4) assume a lossless line ($R' = G' = 0$). For a superconducting CPW operated well below T_c , this is an excellent approximation: the series resistance R' arising from quasiparticle absorption is exponentially suppressed as $\exp(-\Delta/k_B T)$ [53, 54], and the lossless limit is valid for $T \lesssim T_c/2$.

A resonator is formed by taking a finite section of transmission line of length ℓ and imposing boundary conditions at its two ends. These boundary conditions select only those wavelengths that are commensurate with the cavity, defining a discrete set of resonance frequencies. In the CPW resonators of this work, two geometries are used (Fig. 2.2):

Half-wavelength ($\lambda/2$) resonators. Both ends are coupled to the microwave circuit through small gaps that form coupling capacitors. At resonance, they impose current nodes at both ends $I(0) = I(\ell) = 0$. The resulting standing-wave condition $\ell = n\lambda/2$ yields resonance frequencies

$$f_n = \frac{n v_{\text{ph}}}{2\ell} = \frac{n}{2\ell\sqrt{L'C'}}, \quad n = 1, 2, 3, \dots \quad (2.5)$$

The fundamental ($n = 1$) current profile is $I(x) = I_0 \sin(\pi x/\ell)$, with a single current antinode at the center and current nodes at both ends. Harmonics have n antinodes and are equally spaced in frequency.

Quarter-wavelength ($\lambda/4$) resonators. One end is short-circuited to the ground plane (imposing $V(0) = 0$ and a current antinode), while the other end is coupled to the feedline through a capacitor (imposing a current node $I(\ell) = 0$). The resonance condition $\ell = (2n - 1)\lambda/4$ gives

$$f_n = \frac{(2n - 1) v_{\text{ph}}}{4\ell}, \quad n = 1, 2, 3, \dots \quad (2.6)$$

The fundamental current profile is $I(x) = I_0 \cos(\pi x/(2\ell))$, with maximum current at the short. This geometry produces only odd harmonics.

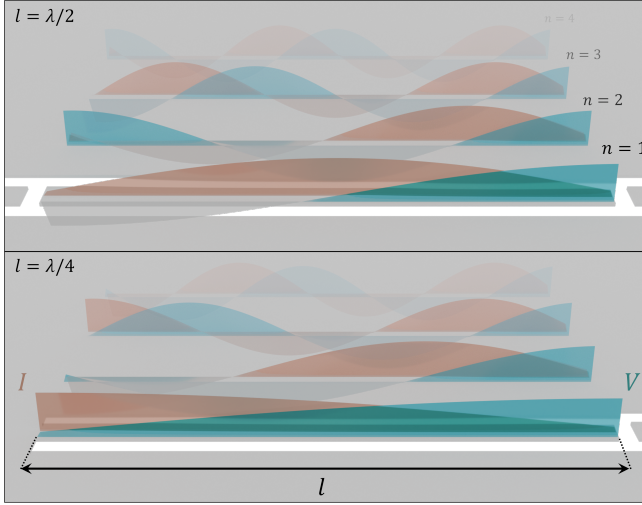


Figure 2.2: Standing-wave current $I(x)$ (orange) and voltage $V(x)$ (blue) profiles for the fundamental mode of (*top*) a $\lambda/2$ resonator, with current nodes at both ends and a current antinode at the center; and (*bottom*) a $\lambda/4$ resonator, with a current antinode at the shorted end and a current node at the coupled end. Equations (2.5) and (2.6) give the corresponding resonance frequencies. The first harmonics are also represented.

An important subtlety is that the current density is not uniform even within the cross-section of the central conductor: it peaks strongly at the edges of the track and decays toward the center [55]. This lateral non-uniformity, which is a consequence of the electromagnetic boundary conditions at the superconductor surface, has the practical implication that the edges of the resonator are the most sensitive regions to local perturbations (for example vortex penetration or thermal hotspots).

2.1.2 Distributed parameters: geometry and substrate effects

The phase velocity v_{ph} and characteristic impedance Z_0 of the CPW are completely determined by the inductance L' and capacitance C' per unit length, which in turn depend only on the conductor geometry and the electric permittivity of the surrounding media. For the CPW geometry, both quantities admit exact analytical expressions in terms of complete elliptic integrals, derived via conformal mapping [49].

Box 2.2 – CPW distributed parameters

A CPW with central conductor width w and gap width g on a substrate of dielectric constant ε_r and thickness h defines the geometric modulus and its complement [49, 50]:

$$k_0 \equiv \frac{w}{w + 2g}, \quad k'_0 \equiv \sqrt{1 - k_0^2}. \quad (2.7)$$

In the limit of a thick substrate ($h \gg w + 2g$), the electric field is split equally, in terms of geometry, between the air half-space (above) and the substrate half-space (below). The effective permittivity is then [49]:

$$\varepsilon_{\text{eff}} = \frac{1 + \varepsilon_r}{2}. \quad (2.8)$$

For a substrate of finite thickness, ε_{eff} involves a second elliptic-integral ratio evaluated at $k_1 = \sinh(\pi w/4h)/\sinh(\pi(w + 2g)/4h)$ [50]; in the experiments of this thesis, $h \gg w + 2g$ in all cases, so Eq. (2.8) holds. The geometric inductance L'_g and capacitance per unit length C' are [49, 50]:

$$L'_g = \frac{\mu_0}{4} \frac{K(k'_0)}{K(k_0)}, \quad C' = 4\varepsilon_0\varepsilon_{\text{eff}} \frac{K(k_0)}{K(k'_0)}. \quad (2.9)$$

where the complete elliptic integral of the first kind is given by

$$K(k) = \int_0^{\pi/2} \frac{d\theta}{\sqrt{1 - k^2 \sin^2 \theta}}. \quad (2.10)$$

The characteristic impedance and phase velocity follow directly:

$$Z_0 = \sqrt{\frac{L'_g}{C'}} = \frac{Z_{\text{vac}}}{4\sqrt{\varepsilon_{\text{eff}}}} \frac{K(k'_0)}{K(k_0)}, \quad v_{\text{ph}} = \frac{c}{\sqrt{\varepsilon_{\text{eff}}}}, \quad (2.11)$$

where $Z_{\text{vac}} = \sqrt{\mu_0/\varepsilon_0} \approx 377 \, \Omega$. The ratio $K(k'_0)/K(k_0)$ is a monotonically decreasing function of k_0 : widening the central conductor relative to the gap lowers the impedance. For a target $Z_0 = 50 \, \Omega$, the required aspect ratio w/g is uniquely determined by ε_r alone; typical values are $w/g \approx 2$ for sapphire ($\varepsilon_r \approx 9.3$ – 11.5) and $w/g \approx 0.4$ for $(\text{La}_{0.3}\text{Sr}_{0.7})(\text{Al}_{0.65}\text{Ta}_{0.35})\text{O}_3$ (LSAT, $\varepsilon_r \approx 22$).

Beyond setting v_{ph} and Z_0 , the substrate has a direct impact on the resonator's energy losses. The dielectric loss tangent $\tan \delta = \varepsilon''/\varepsilon'$ of the substrate material contributes to the internal quality factor through the dielectric participation ratio p , defined as the fraction of the total electric field energy stored in the substrate [56, 57]:

$$\frac{1}{Q_{\text{diel}}} = p \tan \delta. \quad (2.12)$$

For a thick substrate and the lossless-line geometry of Box 2.2, the participation ratio is given by [50, 58]:

$$p = \frac{\varepsilon_r}{1 + \varepsilon_r}. \quad (2.13)$$

For $\varepsilon_r \gg 1$, $p \rightarrow 1$ and nearly all the electric field energy resides in the substrate, making the loss tangent of the substrate material

the dominant dielectric loss channel. The substrates used in this thesis are selected to minimize $\tan \delta$ while maintaining structural and thermal compatibility with the superconducting film: LSAT ($\epsilon_r \approx 22$, $\tan \delta < 10^{-4}$ at low temperature) for the YBCO resonators [59] and C-plane sapphire (Al_2O_3 , $\epsilon_r \approx 9.3\text{--}11.5$ depending on crystal orientation, $\tan \delta < 10^{-7}$ at low temperature) for the NbTiN resonators [42].

2.1.3 Microwave characterization

Two coupling geometries are used in this work, each producing a distinct S_{21} (see Box 2.3 for a general definition) lineshape. The NbTiN $\lambda/4$ resonators are in the *hanger* (notch-filter) geometry: the resonator is coupled to a through-feedline via a single capacitor, and the resonance appears as a dip against a non-zero background. The YBCO $\lambda/2$ resonators are in the *transmission* (bandpass) geometry: the feedline is interrupted at two points and the resonator bridges the two gaps, so that $S_{21} \approx 0$ off resonance and the resonance appears as a Lorentzian transmission peak. Therefore, no electrical DC measurement of T_c is possible through the feedline in the transmission geometry.

Box 2.3 – Scattering parameters for a two-port network

The scattering matrix (S -matrix) is the standard framework for characterizing linear microwave networks [51]. For a two-port network with characteristic impedance Z_0 , the S -parameters relate incoming wave amplitudes a_i and outgoing wave amplitudes b_i (each normalized to $\sqrt{P_{\text{in}}} \propto V_{\text{in}}$, with P_{in} the input power and V_{in} the input voltage):

$$\begin{pmatrix} b_1 \\ b_2 \end{pmatrix} = \begin{pmatrix} S_{11} & S_{12} \\ S_{21} & S_{22} \end{pmatrix} \begin{pmatrix} a_1 \\ a_2 \end{pmatrix}. \quad (2.14)$$

The diagonal elements S_{11} and S_{22} are the complex reflection coefficients at each port; the off-diagonal elements S_{21} and S_{12}

are the complex transmission coefficients. If port 1 acts as a source, the transmitted power fraction is $|S_{21}|^2$, and the reflected power fraction is $|S_{11}|^2$; for a passive, lossless, reciprocal network, energy conservation requires $|S_{11}|^2 + |S_{21}|^2 = 1$. In practice, $|S_{21}|$ is displayed in decibels: $20 \log_{10} |S_{21}|$, where 0 dB corresponds to full transmission and a large negative value to a deep resonance dip or strong absorption. A VNA measures both the magnitude and phase of the S -parameters as a function of frequency, enabling complete characterization of the network.

In the vicinity of f_0 , both lineshapes are described by the same fitting function [60, 61]:

$$S_{21}(f) = \left(\underbrace{1}_{\text{hanger}} - \frac{Q_L |\hat{Q}_e^{-1}| e^{i\varphi}}{1 + 2iQ_L \frac{f - f_0}{f_0}} \right) (Af + B), \quad (2.15)$$

where Q_L is the total (loaded) quality factor, $\hat{Q}_e = |Q_e|e^{i\varphi}$ is a complex coupling quality factor accounting for impedance mismatch at the coupling point (Box 2.4), and $Af + B$ is a slowly varying linear background. For the hanger geometry, the resonance appears as a dip of depth Q_L/Q_e (see Fig. 2.3). For the transmission geometry the Lorentzian term only survives at resonance and the resonance appears as a peak of height Q_L/Q_e . In both cases the fit yields f_0 , Q_L , and $|Q_e|$, from which Q_i can be recovered as described in Box 2.4.

Box 2.4 – Quality factors of a coupled resonator

The quality factor of a resonator is defined as [51]:

$$Q \equiv \omega_0 \frac{\langle U \rangle}{P_{\text{diss}}} = \frac{f_0}{\Delta f_{-3\text{dB}}}, \quad (2.16)$$

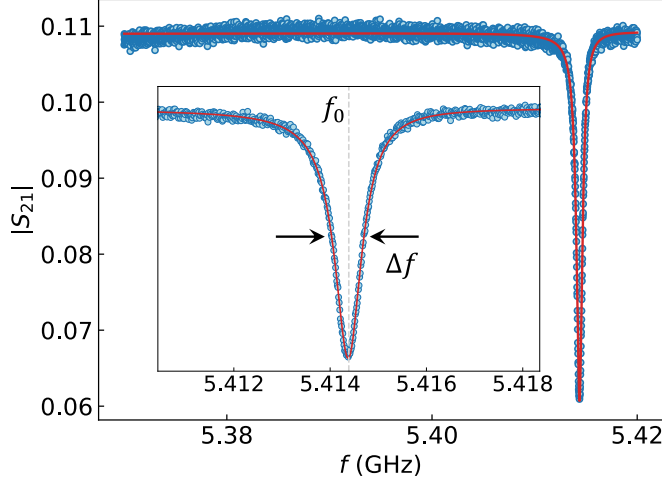


Figure 2.3: Transmission coefficient $|S_{21}(f)|$ of a NbTiN CPW resonator in the hanger geometry, showing the Lorentzian dip at f_0 . The red curve is the fit of Eq. 2.15 on the data. The full width at half maximum $\Delta f = f_0/Q_L$ is indicated by the arrows.

where $\langle U \rangle$ is the time-averaged stored energy, P_{diss} is the dissipated power, and $\Delta f_{-3\text{dB}}$ is the full width at half maximum of the power response. Three quality factors characterize a resonator:

- **Internal quality factor Q_i :** encodes all energy loss mechanisms intrinsic to the resonator itself. Its inverse is additive over all dissipative channels: $Q_i^{-1} = Q_{\text{diel}}^{-1} + Q_{\text{rad}}^{-1} + Q_{\text{qp}}^{-1} + Q_{\text{vortex}}^{-1} + \dots$, accounting for dielectric loss (Eq. (2.12)), radiation, quasiparticle absorption, and vortex-induced dissipation [62], respectively.
- **External (coupling) quality factor Q_e :** quantifies the energy coupled out of the resonator into the feedline per cycle. It is set by the coupling capacitor geometry and is independent of material losses.

- **Loaded quality factor Q_L :** the directly measured FWHM quality factor. In the absence of impedance mismatch ($\varphi = 0$):

$$\frac{1}{Q_L} = \frac{1}{Q_i} + \frac{1}{Q_e}. \quad (2.17)$$

Three coupling regimes are distinguished: *under-coupled* ($Q_e \gg Q_i$, dominated by internal loss), *critically coupled* ($Q_e = Q_i$), and *over-coupled* ($Q_e \ll Q_i$, dominated by coupling loss, $Q_L \approx Q_e$). The resonators in this thesis are deliberately designed in the under-coupled regime ($Q_e \gg Q_i$, so $Q_L \approx Q_i$) [28, 47], ensuring that changes in the material properties (flux penetration, photodoping) are reflected directly in Q_L .

2.1.4 Kinetic inductance: a uniquely superconducting contribution

The expressions for L'_g and C' given in Box 2.2 treat the inductance as purely of geometric (magnetic) origin, which is a property of the field distribution around the conductor and is independent of its material. In a superconductor, however, a second inductive contribution arises from the quantum condensate of Cooper pairs. Cooper pairs have a finite effective mass $m^* = 2m_e$, and when a current flows they carry a kinetic energy density $u_{\text{kin}} = \frac{1}{2}m^*n_s v_s^2$, whose volume integral U_{kin} is proportional to the square of the current (Box 2.5). Since $U_{\text{kin}} = \frac{1}{2}L_k I^2$ by definition of an inductance, this kinetic energy behaves as an additional inductive element in series with the geometric inductance [53, 63]. The resulting contribution is called the *kinetic inductance* L_k . An inductance can also be interpreted as a delay between the applied voltage and the circulating current. In the present case, the kinetic inductance represents the delay caused by the inertia of massive Cooper pairs to react to an oscillating voltage.

Box 2.5 – Origin of kinetic inductance

The electrodynamics of the superconducting condensate are governed by the *first London equation* [63]:

$$\frac{\partial \mathbf{J}_s}{\partial t} = \frac{n_s(2e)^2}{m^*} \mathbf{E} = \frac{\mathbf{E}}{\mu_0 \lambda_L^2}, \quad (2.18)$$

where $\mathbf{J}_s = n_s(2e)\mathbf{v}_s$ is the supercurrent density, n_s is the Cooper-pair density (counting pairs with charge $2e$ and mass $m^* = 2m_e$), and

$$\lambda_L = \sqrt{\frac{m^*}{\mu_0 n_s (2e)^2}} \quad (2.19)$$

is the *London penetration depth*. Equation (2.18) is purely inductive: $\partial \mathbf{J}_s / \partial t \propto \mathbf{E}$ with no resistive component. For an AC excitation at angular frequency ω , it gives a purely imaginary surface impedance $Z_s = i\omega\mu_0\lambda_L^2$, confirming the inductive character. For a superconducting strip of length ℓ , width w , and film thickness $d \ll \lambda_L$ (the thin-film limit), the kinetic energy of the superfluid expressed in terms of the total current I_s is:

$$U_{\text{kin}} = \frac{1}{2} m^* n_s \int v_s^2 dV = \frac{m^* \ell}{2 n_s (2e)^2 w d} I_s^2 \equiv \frac{1}{2} L_k I_s^2, \quad (2.20)$$

which identifies the kinetic inductance per unit length that simply adds up to the geometric one:

$$L'_k = \frac{\mu_0 \lambda_L^2}{w d}. \quad (2.21)$$

For a CPW geometry, the non-uniform current distribution across the strip transverse direction introduces a geometric correction factor $g_k(w, g)$ of order unity [58, 64]. An equivalent expression in terms of the normal-state sheet resistance $R_\square$ and the superconducting gap $\Delta(T)$, derived from Mattis–Bardeen theory [54],

is:

$$L'_k(T) = \frac{\hbar R_\square}{\pi w \Delta(T)} \frac{1}{\tanh(\Delta(T)/2k_B T)}. \quad (2.22)$$

Both forms show that $L'_k \propto \lambda_L^2 \propto 1/n_s$: *the kinetic inductance is inversely proportional to the superfluid density, and therefore diverges at T_c .*

The total inductance per unit length of a superconducting CPW is $L' = L'_g + L'_k$, and the resonance frequency accordingly becomes:

$$f_0 = \frac{\alpha}{2\ell \sqrt{(L'_g + L'_k) C'}}, \quad (2.23)$$

with $\alpha = 1$ (resp. $1/2$) for a $\lambda/2$ (resp. $\lambda/4$) resonator. The kinetic inductance fraction $\alpha_k \equiv L'_k/(L'_g + L'_k)$ quantifies the relative weight of this contribution. For both the NbTiN resonators of Chapter 5 and the YBCO resonators of Chapter 6, the kinetic part of the total inductance remains marginal, of the order of 10%.

Crucially, $L'_k \propto 1/n_s$ and n_s is a function of temperature, so the resonance frequency f_0 is intrinsically temperature-dependent. As T increases toward T_c , n_s decreases toward zero: fewer and fewer Cooper pairs are available. Box 2.6 gives the quantitative form of this dependence. Since the total supercurrent must be maintained, each Cooper pair must carry a larger share as their number dwindles, implying a larger drift velocity $v_s \propto I_s/(n_s(2e))$. More kinetic energy is therefore stored per unit current squared, which is precisely what a larger inductance means. The kinetic inductance thus diverges as $T \rightarrow T_c$, driving f_0 toward zero.

Box 2.6 – Two-fluid model and the superfluid density

The two-fluid model, proposed phenomenologically by Gorter and Casimir in 1934 [65], and the microscopic justification by BCS theory [53], partitions the electron system into a superfluid component and a normal-fluid component of density $n_n(T)$, with $2n_s(T) + n_n(T) = n$, where n is the total electron density and n_s remains the Cooper-pair density of Box 2.5 (each pair removing two electrons from the normal fluid). The model is often written in terms of the condensed *electron* density $2n_s$; since the two differ by a constant factor, the temperature dependence and the scaling $L'_k \propto 1/n_s$ are identical in either counting. In the Gorter–Casimir model:

$$\frac{2n_s(T)}{n} = 1 - \left(\frac{T}{T_c}\right)^4, \quad (2.24)$$

so that $n_s(T = 0) = n/2$ and $n_s(T = T_c) = 0$. This exponent 4 applies to conventional isotropic (three-dimensional, *s*-wave) superconductors. For the high- T_c cuprate YBCO used in Chapter 6, the *d*-wave symmetry of the order parameter and the quasi-two-dimensional electronic structure lead to deviations [53]; the temperature dependence of L'_k is instead fitted to a generalized form of Eq. (2.24), in which the exponent 4 is replaced by a free exponent a , so that $L'_k \propto 1/n_s$ gives:

$$L'_k(T) = \frac{L'_{k,0}}{1 - (T/T_c)^a}, \quad (2.25)$$

with the exponent a treated as a fitting parameter (typically $a \approx 1.5$ – 2 for YBCO [47, 66]). For conventional NbTiN, the full Mattis–Bardeen expression (Eq. (2.22), Box 2.5) can be used instead, with the gap approximated by [53]:

$$\Delta(T) \approx \Delta(0) \tanh\left(1.74\sqrt{\frac{T_c}{T} - 1}\right). \quad (2.26)$$

In both cases, as $T \rightarrow T_c$: either $n_s \rightarrow 0$ (Eq. (2.24)) or $\Delta \rightarrow 0$ (Eq. (2.26)), and $L'_k(T) \rightarrow \infty$.

The divergence of L'_k causes the resonance frequency to decrease steeply near T_c , following Eq. (2.23). The resulting $f_0(T)$ curve displays an asymptotic plateau at $T \ll T_c$ giving way to a sharp downward curvature near T_c . Some examples measured on NbTiN and YBCO resonators are shown in Fig. 2.4 where both fitting strategies (fitting Eq. (2.23) with either Eq. (2.25) or (2.22)) are compared. This standard procedure yields the material parameters $L'_{k,0}$, T_c , and $R_\square$. The deeper consequence of these expressions is that f_0 is sensitive to any perturbation of n_s , whether driven by temperature (uniform, or local as in Chapter 6), penetrating vortices, or optical Cooper-pair breaking. This single fact that the resonance frequency transduces changes in the superfluid density underpins the two experimental chapters that follow.

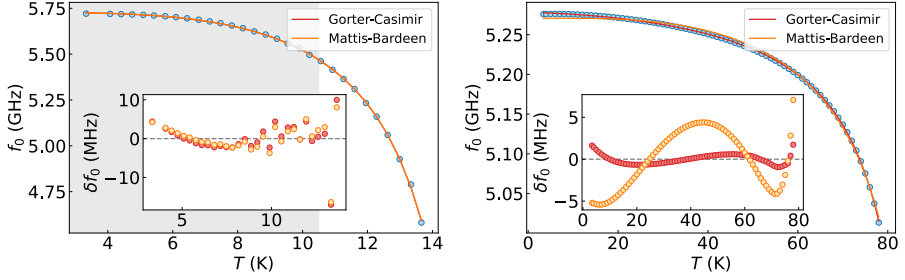


Figure 2.4: (*Left*) Resonance frequency $f_0(T)$ of a NbTiN superconducting CPW resonator, showing the plateau at $T \ll T_c$ and the steep decrease toward T_c driven by the divergence of L'_k . Lines indicate the Gorter–Casimir (Eq. (2.25), the fitting exponent yields 3.76 ± 0.11) and Mattis–Bardeen (Eq. (2.22)) fits. The inset shows the fitting error. The shaded region ($T < T_c/2$) marks the temperature range in which flux avalanche activity is observed in NbTiN (Chapter 5). (*Right*) Resonance frequency $f_0(T)$ of an optimally doped YBCO superconducting CPW resonator. The Mattis–Bardeen fit fails to reproduce the weaker plateau at $T \ll T_c$ because YBCO is not described by BCS theory. The Gorter–Casimir form captures this discrepancy and yields a fitting exponent of 1.60 ± 0.03 .

Chapter 3

Magneto-Optical Imaging: Principle and Practices

Mapping the magnetic field distribution at the surface of a superconductor is one of the most direct routes to assess its electromagnetic response. The spatial pattern of the out-of-plane flux density $B_z(x, y)$ encodes the location of vortices, the topology of screening currents, and the influence of material inhomogeneities, information that bulk measurements inherently average away. This chapter describes the magneto-optical imaging (MOI) technique, which provides this information with micrometer spatial resolution over millimeter-scale fields of view, and underpins the experimental work reported in Chapters 5 and 7.

3.1 Physical principle

Several local techniques exist for imaging B_z in a superconductor [67]. Scanning Hall probe microscopy (SHPM) rasters a sub-millimeter Hall sensor across the sample and provides a quantitative, calibrated field map [68, 69]; however, its point-by-point acquisition is intrinsically slow and requires a mechanically stable cryogenic scanning stage. Scanning SQUID microscopy achieves field sensitivities well below a

nanotesla [68] at comparable mechanical complexity, while magnetic force microscopy (MFM) offers nanometer-scale resolution but at the cost of a potentially invasive probe tip that can perturb the vortex state being imaged [70]. Magneto-optical imaging circumvents these limitations by converting the entire $B_z(x, y)$ distribution into a light-intensity map captured in a single widefield camera frame, combining a large field of view ($\sim \text{mm}^2$), sub-second acquisition, and little magnetic invasiveness [71–73].

The operating principle of MOI relies on the Faraday effect in a magneto-optically active indicator film placed on top of the sample. The indicator possesses a large Verdet constant and a spontaneous magnetization fixed in plane by the Bi doping. A local field B_z tilts this magnetization slightly out of plane, rotating the polarization of a linearly polarized beam that traverses the indicator proportionally to the magnetization component along the light path. After reflection from an aluminum mirror at the bottom of the indicator stack, the beam makes a second pass through the active layer (doubling the effective rotation) before passing through an analyzer whose transmission axis is nearly crossed with the initial polarizer. The resulting intensity distribution, recorded by a CCD or a CMOS camera, constitutes a pixel-wise map of B_z at the indicator plane. The physical mechanisms governing this conversion are formalized in Box 3.1.

Box 3.1 – The Faraday effect in a Bi:YIG indicator

The indicator used throughout this work is a Bismuth-doped Yttrium Iron Garnet film ($\text{Y}_{3-x}\text{Bi}_x\text{Fe}_5\text{O}_{12}$, Bi:YIG) of thickness $d \simeq 3 \mu\text{m}$, epitaxially grown on a $450 \mu\text{m}$ -thick Gadolinium Gallium Garnet (GGG, $\text{Gd}_3\text{Ga}_5\text{O}_{12}$) substrate. The GGG provides a matched lattice parameter for epitaxial Bi:YIG growth and is optically transparent at the operating wavelength. A 100 nm aluminum mirror is deposited on the optically active face of the indicator.

Faraday rotation. When linearly polarized light propagates along

$\hat{z}$ through a medium of magnetization $\mathbf{M}$, the polarization plane rotates by the Faraday angle [74]

$$\theta_F = V(\lambda) \mu_0 \int_{\text{o.p.}} \mathbf{M} \cdot d\mathbf{l}, \quad (3.1)$$

where $V(\lambda)$ is the wavelength-dependent Verdet constant and the integral is taken along the optical path. For normal incidence on a homogeneous film of thickness d , and accounting for the double pass upon reflection by the Al mirror,

$$\theta_F = 2 V(\lambda) \mu_0 M_z d. \quad (3.2)$$

At $\lambda = 550 \text{ nm}$, the Verdet constant of the Bi:YIG indicators used in this thesis is $V_{\text{Bi:YIG}} = 0.018 \pm 0.005 \text{ deg mT}^{-1} \mu\text{m}^{-1}$ [74].

Indicator response to B_z . Due to its magnetocrystalline anisotropy induced by the Bi doping, the saturation magnetization $\mathbf{M}_s$ lies in the film plane at zero field. A perpendicular field B_z tilts $\mathbf{M}_s$ by an angle ϕ from the plane. Minimizing the magnetic energy

$$E_{\text{mag}} = E_a \sin^2 \phi - B_z M_s \sin \phi \quad (3.3)$$

yields the equilibrium condition

$$\tan \phi = \frac{B_z}{B_s}, \quad B_s \equiv \frac{2E_a}{M_s}, \quad (3.4)$$

where $B_s \approx 80 \text{ mT}$ is the anisotropy saturation field of the indicator and E_a is the uniaxial anisotropy energy density [75]. Substituting $M_z = M_s \sin \phi$ into Eq. (3.2),

$$\theta_F = C M_s \sin \left(\arctan \frac{B_z}{B_s} \right), \quad C \equiv 2V(\lambda) \mu_0 d. \quad (3.5)$$

The response is linear for $B_z \ll B_s$ and saturates near $B_z \sim B_s$.

Detected intensity. After the analyzer, set at a small angle $\beta(x, y)$ from the 90° extinction, Malus' law gives

$$I(B_z, x, y) = I_0(x, y) \sin^2[\theta_F(B_z, x, y) + \beta(x, y)] . \quad (3.6)$$

The spatial variation of $\beta(x, y)$ arises from the finite numerical aperture of the objective and any inclination of the indicator relative to the applied field. It is accounted for pixel by pixel during calibration (see Section 3.3.1).

3.2 Experimental setup and cryogenic environment

The polarization microscope is a modular commercial system (Olympus BX-RLA2). A 56.4 W LED lamp (CoolLED pE-300 white) provides the broadband source, filtered to $\lambda = 550$ nm by a bandpass element (Olympus U-25IF550); this wavelength is chosen to coincide with the emission peak of the LED and with the region of highest Verdet constant for Bi:YIG indicators [74]. The filtered beam passes through a fixed linear polarizer, then is deflected by a beam splitter toward the sample through a $\times 5$ objective (LMPLFLN 5 BD), and impinges on the indicator. After the double pass through the Bi:YIG active layer and reflection from the Al mirror, the returning beam is transmitted through the beam splitter to an analyzer (U-AN360-3, angular precision 0.1°). The signal is finally recorded by a 2048×2048 pixel, 12-bit CCD camera (RETIGA-4000R). A schematic of the full optical path is shown in Fig. 3.1.

The sample is maintained at cryogenic temperature in a closed-cycle helium cryostat (Montana Instruments Cryostation C1) with a base temperature of ~ 3.2 K. An external cylindrical copper solenoid (resistance $R = 22 \, \Omega$, inductance $L = 26$ mH), fed by a precision current source (Keithley 2400 series), generates a perpendicular applied

field with spatial homogeneity better than 1% over a $5 \times 5 \text{ mm}^2$ area at the sample location [75]. The entire assembly rests on an actively damped, non-magnetic optical table.

The Bi:YIG indicator is placed face-down on the sample (*i.e.* the active layer is as close as possible to the sample surface). The air gap between the Al mirror and the sample surface, set by surface roughness, dust particles, and mounting conditions, is typically $10 \mu\text{m}$ for standard free-standing placement, and constitutes the dominant factor limiting spatial resolution, which is of the same order. The gap can be reduced to $\sim 2 \mu\text{m}$ by pressing the indicator against the sample with a mechanical clip, improving spatial resolution at the cost of greater mechanical stress on the indicator and a higher propensity for magnetic domain pollution.

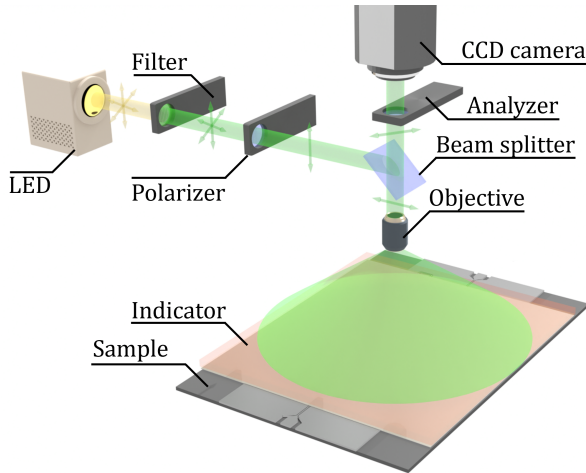


Figure 3.1: Schematic of the magneto-optical imaging setup. A filtered LED lamp provides monochromatic 550 nm illumination that is polarized, reflected onto the indicator-covered sample inside a cryostat, and collected by a nearly-crossed analyzer before detection by a CCD camera. This schematic representation is not to scale, as the sample is usually much smaller than the objective.

Incompatibility with radio-frequency measurements. The 100 nm Al reflective coating on the indicator is essential for achieving a comfortable signal-to-noise ratio by doubling the optical path through the Bi:YIG layer and by providing a nearly uniform background irrespective of the topography and design of the sample. On top of that, it avoids the interaction of light with the sample. At GHz frequencies, however, this conducting layer supports eddy currents driven by the microwave field, strongly attenuating the RF signal inside a superconducting resonator placed beneath it. *In situ* combination of MOI with microwave spectroscopy is therefore not feasible with this indicator; the experimental consequences for the NbTiN resonator measurements are discussed in Chapter 5.

Magnetic domain artifacts. The Bi:YIG layer is subdivided into in-plane magnetic domains with different orientations of $\mathbf{M}_s$. When the indicator is slightly tilted with respect to the applied field direction, neighboring domains exhibit different Faraday rotations (Eq. (3.5)), producing bright or dark triangular patches polluting the raw image and rendering the post-processing very challenging [74, 75]. Their impact can be reduced by applying a small in-plane bias field (~ 2 mT) to force a single-domain state in the indicator (though at the cost of reduced perpendicular-field sensitivity).

3.3 Post-processing

3.3.1 Calibration: intensity to magnetic field

Raw CCD images carry intensity values governed by Eq. (3.6), but the three spatially varying parameters, incident intensity $I_0(x, y)$, local sensitivity $A(x, y)$, and analyzer offset $\beta(x, y)$, must be determined empirically for each experimental configuration. In the regime $B_z \ll B_s$, the Taylor expansion of Eq. (3.6) reduces to a parabola in B_z [75]:

$$I(B_z, x, y) \approx I_{\min}(x, y) + A(x, y) [B_z(x, y) - B_{\min}(x, y)]^2. \quad (3.7)$$

Here $B_{\min}(x, y)$ is the local field at which the intensity reaches its minimum $I_{\min}(x, y)$, and $A(x, y)$ encodes the local sensitivity of the indicator.

The three parameter maps are extracted from a calibration dataset acquired at $T \gtrsim T_c$, where the superconducting film is in the normal state and the field at every pixel equals the uniform applied field $\mu_0 H$. A common measurement protocol consists of acquiring sequence of images recorded while sweeping H from -4 mT to $+4$ mT in steps of 0.1 mT, with several frames averaged at each step to reduce shot noise. The intensity profile $I(H)$ at each pixel (x, y) is fitted to Eq. (3.7), yielding the pixel-wise calibration maps $\{I_{\min}, B_{\min}, A\}(x, y)$.

Any subsequent image acquired below T_c is converted to $B_z(x, y)$ by inverting Eq. (3.7) pixel by pixel by using the acquired calibration maps. The analyzer is deliberately set $\sim 4^\circ$ from maximum extinction so that the intensity minimum falls outside the measurement field range, ensuring that each intensity value unambiguously corresponds to a single branch of the calibration parabola. The field resolution achieved by this procedure is ≈ 0.1 mT for a point-like source, improving to ≈ 0.01 mT for spatially extended sources where pixel averaging can be exploited.

3.3.2 MOI inversion: from flux density to current density

The calibrated $B_z(x, y)$ map gives the perpendicular flux density at the indicator plane, located a height $h \approx 10$ μm above the superconducting film. The in-plane supercurrent density $\mathbf{J}(x, y)$ responsible for this field is related to B_z by the non-local Biot–Savart law. Recovering $\mathbf{J}$ requires a full-field inversion that cannot be reduced to a local operation. The computationally efficient FFT-based scheme introduced by Jooss *et al.* [76], and employed in the experimental analysis of Chapter 7, performs this inversion with a computational cost that scales as $N \log N$ rather than N^3 for an N -pixel image. The key mathematical steps are summarized in Box 3.2.

Box 3.2 – Biot–Savart inversion for thin-film

For a superconducting film thinner than the London penetration depth, the supercurrent density $\mathbf{J}(x, y)$ lies in the film plane and satisfies $\nabla \cdot \mathbf{J} = 0$. It is conveniently expressed through a scalar stream function $g(x, y)$:

$$\mathbf{J} = \nabla \times \hat{z} g. \quad (3.8)$$

The Biot–Savart relation between g and the measured B_z becomes a product in Fourier space:

$$\tilde{B}_z(\mathbf{k}) = \mu_0 \tilde{K}_g(k, h, d_s) \tilde{g}(\mathbf{k}), \quad (3.9)$$

with the analytical kernel ($k = |\mathbf{k}|$, measurement height h , superconducting film thickness d_s)

$$\tilde{K}_g(k, h, d_s) = e^{-kh} \sinh\left(\frac{kd_s}{2}\right). \quad (3.10)$$

Inverting Eq. (3.9) gives $\tilde{g}$; the current components then follow as $\tilde{J}_x = -ik_y \tilde{g}$ and $\tilde{J}_y = ik_x \tilde{g}$.

In practice, $\tilde{K}_g^{-1}$ acts as a high-pass filter that amplifies measurement noise at large k ; a Hanning window is applied in Fourier space to regularize the inversion. To suppress aliasing artifacts from the periodic continuation implicit in the FFT, the measured image is padded with the value of the applied magnetic field to approximately three times the sample area before transformation [77]. The indicator-to-sample distance $h \simeq 10 \mu\text{m}$ and the superconducting film thickness d_s (200 nm for the Nb films of Chapter 7) enter as fixed parameters of the kernel (3.10). Note that d_s is the thickness of the superconducting film under study, not that of the Bi:YIG indicator, which enters only through the standoff distance h .

The current density streamlines reconstructed in this way, super-

imposed on calibrated B_z maps, are used in Chapter 7 to identify the role of LIPSS domains in redirecting the screening current flow and to validate critical-state model simulations.

Chapter 4

Scanning Laser Microscopy

Scanning laser microscopy (SLM) is a technique that raster-scans the surface of a sample under test (SUT) pixel by pixel using a focused laser beam as the optical probe. The interaction between light and the SUT locally alters the properties of the sample (or of the probe) and gives rise to a position-dependent photoresponse, $\text{Pr}(x, y)$.

The technique has reached maturity and is routinely applied across a broad range of disciplines, including biology [78, 79], engineering [80], and materials science [81]. Several books [82] and review articles [83–86] have examined its fundamentals and practical realization, and have proposed robust instrument designs. Commercial confocal microscopes are widely available, but they are built for reflectivity and fluorescence imaging at ambient conditions; probing the microwave response of a superconducting device requires the beam to be delivered into an optical cryostat and the photoresponse to be read out from the device itself, which is why such instruments are, as a rule, assembled in the laboratory.

Depending on the chosen photoresponse, a variety of imaging modes can be implemented, selected according to the SUT and the specific sample property to be probed. The most common is reflectometry, in which $\text{Pr}(x, y)$ is modulated by variations in the local refractive index and surface orientation.

SLM has proven particularly useful in the study of superconducting films, for which integration of the setup in an optical cryostat is essential [38, 87–94]. For a current-carrying superconducting film, a convenient response signal is the change in the voltage drop across the sample arising from the change in resistivity at the illuminated spot. The primary mechanism governing this change in resistivity is the lattice heating by low-energy phonons. In addition, the photon-induced breaking of Cooper pairs, which generates high-energy quasiparticles, may also need to be considered [95]. Imaging modes other than optical and voltage contrast that reveal the superconducting properties of the SUT are described in Ref. [92].

Although the technique itself is mature, no such instrument was available in Liège at the start of this work. The microscope described in this Chapter was therefore built from the ground up over the course of this thesis: the optical architecture was designed, the components selected and assembled, the control and acquisition software written, and the spatial and thermal responses calibrated. The detection modes of Section 4.2, including the time-resolved multi-harmonic scheme of Section 4.3, were developed on this platform.

This Chapter describes that setup, used in the second part of Chapter 6, together with the practical implementation and core elements of such a microscope and the different sample-probing modes available. It is written as a practical guide rather than a purely theoretical chapter, with enough detail, down to the part numbers of Table 4.1, for the instrument to be reproduced elsewhere.

4.1 Experimental setup

The central architecture of the scanning laser microscope is illustrated in Fig. 4.1(a) (optical elements) and Fig. 4.1(b) (integrated assembly). The system comprises a collimated laser source coupled to a dual-axis galvanometer scanner, which deflects the beam toward a scan lens (see Fig. 4.1(c) for a detailed view). The mirror positions are set by a

control voltage applied to the galvo-controller inputs. The scanner and the scan lens are arranged so that the mean pivot point of the mirrors coincides with the back-focal plane of the scan lens; the beam therefore emerges from the scan lens parallel to the primary optical axis of the microscope. It then passes through a tube lens, positioned so that its back focal point coincides with the image point of the scan lens, thereby forming a $4f$ relay system.

Finally, a 50/50 beam splitter reflects the beam toward a microscope objective placed at the focal point of the tube lens¹. The beam enters the objective pupil with a pivot angle that is linearly converted into a lateral displacement of the focused spot on the sample under test (see Fig. 4.1(d)). The light reflected from the sample passes back through the beam splitter and is collected by a photodiode.

The elements composing the core of the microscope (*i.e.* the essential parts needed to steer and focus the beam on the sample, laser source excluded) are listed in Table 4.1. Unless otherwise specified, the part numbers all correspond to the *Thorlabs* catalog.

4.2 Detection modes

A focused laser beam can give rise to a wide variety of light–matter interactions, each of which produces a distinct photoresponse that can be probed. This Section details the modes developed and used in this thesis, beginning with the most conventional and intuitive one, the measurement of the reflectivity.

¹The microscope can also be converted into a nitrogen-vacancy (NV) microscope, although this is not pursued in detail in this thesis: an NV-rich diamond placed on the sample acts as an indicator (as in MOI, see Chapter 3), and the optically detected magnetic resonance (ODMR) signal is collected while the sample is raster-scanned. In that configuration, the signal-to-noise ratio can be greatly improved by replacing the beam splitter with a dichroic mirror (ref. DMLP567R).

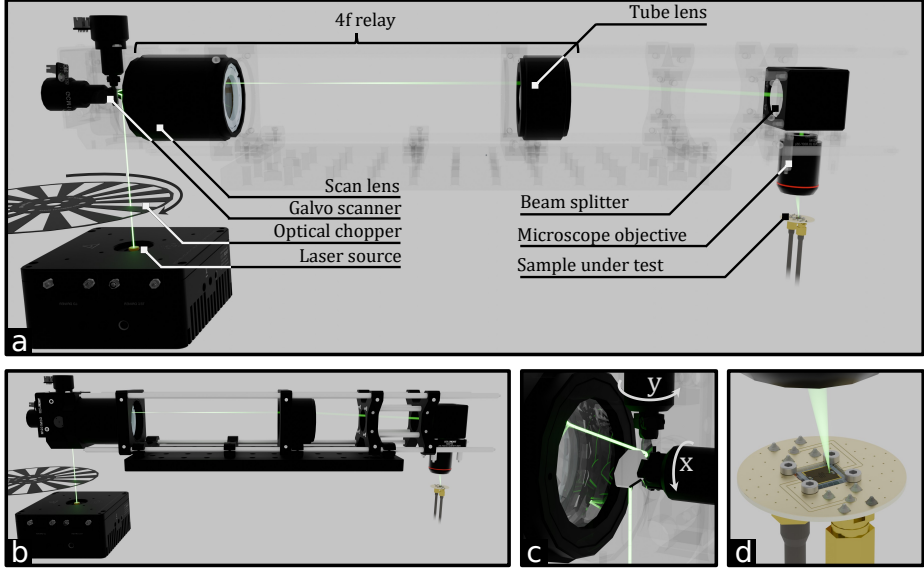


Figure 4.1: (a) Schematic of the optical elements of the scanning laser microscope. The collimated laser beam passes through a mechanical chopper, is steered by a dual-axis galvo scanner, propagates through a $4f$ -relay system, and is reflected toward the objective pupil by a 50/50 beam splitter. A photodiode can be mounted on top of the beam splitter to measure the reflected light intensity. (b) Schematic of the complete mechanical assembly, machined as a single block that can be attached to any three-dimensional positioning mount. (c) Close-up of the dual-axis scanner, whose center is placed at the focal point of the scan lens. (d) Sample under test electrically connected to the instruments used for the susceptibility measurements.

Category	Part number	Purpose
O	GVSK2-EC	Galvanometric scanner (GS)
O	CLS-SL	Scan lens (SL)
O	TTL180-A	Tube lens (TL)
O	CCM1-BS013/M	Beam splitter (BS)
O	<i>Olympus</i> LMPLFLN series	Objective (O)
O	DJ532-40	532 nm laser
S	GCM102/M	Mounting frame for GS
S	SM2L05 + SM2V05	SL mount and focal adjustment
S	LCP34/M	SL and TL cage mounts
S	ER18	Structural rods
S	MB1030/M	Optical breadboard (OB)
S	LCP33B	Cage to OB mounts
S	LCP33/M	Cage adapter for BS
S	SM1QA + SM1A3 + CP33/M	O quick release mount

Table 4.1: Non-exhaustive list of the optical (O) and structural (S) elements used in the SLM setup. All elements are supplied by *Thorlabs*, except for the objectives, which are from *Olympus*.

4.2.1 Optical mode

As shown later, the microscope developed here reaches a resolution down to the sub-micrometer scale. At these scales, the laser spot must be placed reliably and accurately at the desired position on the sample. To this end, a converging lens (LA1951-A-ML) and a single pixel photodiode (SM1PD1A) are mounted on top of the beam splitter and collect the light reflected by the sample at the position where the laser is focused. By raster-scanning a region of interest, the reflectivity is measured pixel-by-pixel and reconstructed into a single image, analogous to the one obtained with a widefield optical microscope. Two distinct notions of resolution should be distinguished here. The *pixel resolution* (the number of pixels in the image) is, in widefield imaging, limited by the camera sensor size (of the order of $10^3 \times 10^3$); in SLM it can be made much larger, since the pixels are acquired sequentially. The *spatial resolution* (the ability to resolve small features) is instead limited by diffraction, and is therefore set

by the Gaussian spot size on the sample.

Determination of the optical spot size

The size of the laser spot on the sample depends on the magnification of the objective used. Furthermore, the X/Y driving voltage-to-distance conversion factor and distortion amplitude are objective-dependent. Therefore, it is necessary to carry out a calibration procedure for each magnification used. To do so, calibration chessboards of different lateral size (each 11×11 squares) made of reflective 50-nm-thick Al were patterned on top of Si/SiO₂ (300 nm) substrates. Fig. 4.2(a) shows an image of one of the chessboard patterns as acquired by the SLM with a $20\times$ objective using $500 \times 500 \text{ px}^2$ resolution. An example of a single 10^3 px line scan along the segment A-B is shown in panel (b). This position-dependent photoresponse is fitted using the expression:

$$\text{Pr}(x, y_0) = A \operatorname{erf}((x - x_0)/\sqrt{2}\sigma) + Bx + C \quad (4.1)$$

with $\operatorname{erf}(x)$ the error function, x_0 the position of the border of the corresponding Al square, σ the standard deviation and A, B, C constants. Fig. 4.2(c) shows the spot size (width of the gaussian intensity distribution at $1/e^2$) of the focused beam, corresponding to four times the standard deviation obtained with Eq. 4.1, for each border of the square on the board. This figure shows that the beam exhibits little distortion across the entire field of view for magnifications below $50\times$, and a tendency to increase in size for $50\times$ and $100\times$ objectives near the field-of-view boundaries. We have also investigated pincushion field distortion. This distortion is caused by the finite distance between the rotation axes of both mirrors and results from the increase of magnification with distance from the optical axis, thereby causing straight lines to bow inward. By using a hyperbolic fit, we found that these distortions have a maximum amplitude of $6 \mu\text{m}$ for a line scan of 1 mm , and can be readily corrected if needed, although they represent a negligible source of error which can typically be ignored. In

Fig. 4.2(d), the dependence of the spot size on the numerical aperture of the objectives is shown. The gray dashed line shows the Airy disk diameter $d = 1.22\lambda/\text{NA}$ (where NA is the numerical aperture of the used objective) expected from the diffraction limit for a laser source of wavelength $\lambda = 532\text{ nm}$. A more rigorous derivation accounting for the gaussian character of the laser yields a smaller diffraction limit, given by $d = 0.635\lambda/\text{NA}$ [96], shown in dashed black.

4.2.2 Electrical modes

In addition to being reflected, a fraction of the incident light is absorbed by the sample and, in the most general case, converted into heat. Consequently, if the transport properties of the sample are temperature-dependent, as is the case for all samples studied in this thesis and in general for most of the electrical samples, raster-scanning the sample reconstructs an image of its local temperature *susceptibility*. The word *susceptibility* is chosen with care: because heat propagation is diffusive, the electrical response acquired when illuminating a single pixel results from the modification of the transport properties across the entire induced temperature variation, which may span a large region. The extent of this region is set by the thermal diffusivity of the sample,

$$D = \frac{\kappa}{\rho C_v}, \quad (4.2)$$

where κ is the thermal conductivity, ρ the mass density and C_v the specific heat of the substrate/sample. The thermal probe can therefore be very local (low diffusivity) or, conversely, very broad (high diffusivity); at best, it is of the same size as the optical probe.

Induced thermal profile

To better characterize the temperature profile generated by light absorption, a 50-nm-wide Nb bridge in a four-wire configuration, fabricated on top of Si/SiO₂ (see inset of Fig. 4.3), is used as a sensitive thermometer. The laser is scanned across the surface while

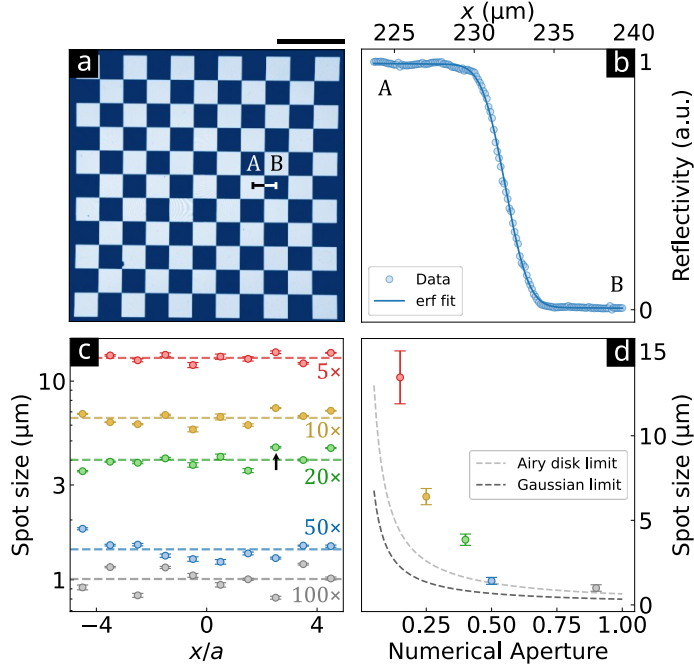


Figure 4.2: (a) Representative $501 \times 501 \text{ px}^2$ SLM image acquired using the sample reflectivity as photoresponse, obtained with a $20\times$ objective and a photodiode as detector. The scale bar in the upper-right corner corresponds to $250 \mu\text{m}$. (b) Line scan along the segment A–B indicated in panel (a), spanning the transition between the substrate (blue) and a reflective Al square (white). The solid line is the fit of Eq. (4.1), which yields the $1/e^2$ beam diameter (spot size). (c) Measured spot size at various x -axis scan positions, normalized by the size a of a single chessboard cell ($a = 272.8, 272.8, 90.9, 22.7$ and $11.4 \mu\text{m}$ for the $5\times, 10\times, 20\times, 50\times$ and $100\times$ objectives, respectively). Dashed lines indicate the mean values, and the black arrow marks the spot size associated with the fit in panel (b). (d) Mean beam diameter for each objective tested. The light-gray dashed line is the Airy-disk limit, spot size $= \frac{1.22 \lambda}{\text{NA}}$, and the dark-gray dashed line is the Gaussian-beam limit, spot size $= \frac{0.635 \lambda}{\text{NA}}$.

the temperature at the junction is monitored through its resistance; tracking this temperature as a function of the spot position amounts to measuring the temperature profile, as represented schematically in Fig. 4.3(b). The resistance of the junction as a function of temperature is shown in Fig. 4.3(a). In the vicinity of the critical temperature T_c (around 8.5 K), the resistance varies by several ohms over a narrow temperature interval, which makes the junction a sensitive local thermometer. More precisely, the technique can resolve the temperature profile only when the laser spot lies in the region where it induces a temperature increase ΔT at the junction such that $T_{\text{base}} + \Delta T \approx T_c$, where T_{base} is the temperature of the cold finger to which the back side of the sample is thermally anchored. When the cold finger is kept at low temperature, ΔT must be large and the profile is measured over small radial distances r (see inset of Fig. 4.3(a)) from the junction; conversely, when the cold finger is held at a temperature close to but below T_c , ΔT can be small and the profile is resolved over large distances.

Following this method, $R(r)$ curves are acquired at various fixed T_{base} and are shown in Fig. 4.3(c). As the laser, fixed at an optical power of 2.5 mW, approaches the junction, a sudden temperature rise is observed once the junction temperature reaches the vicinity of $T_c \approx 8.5$ K. From these curves, the temperature profile $\Delta T(r)$ is reconstructed in Fig. 4.3(d). The profile follows a $1/r$ law, as expected for a steady point source (see Box 4.1, Eq. (4.7)) [97]. The method cannot resolve the temperature in the region close to the optical spot because, as visible in Fig. 4.3(c), even at the lowest cold finger temperature the junction reaches the normal state at a distance of approximately 5 μm . Nevertheless, a clear bump is observed when the laser is close enough to the junction. This rise in resistance occurs in the normal state, where the slope of $R(T)$ is much weaker, so the heating induced by direct illumination is expected to be significantly larger. This is further enhanced by the fact that the heat is deposited directly into the Nb, rather than in the substrate, from which it would otherwise have to cross the substrate–device interface, associated with

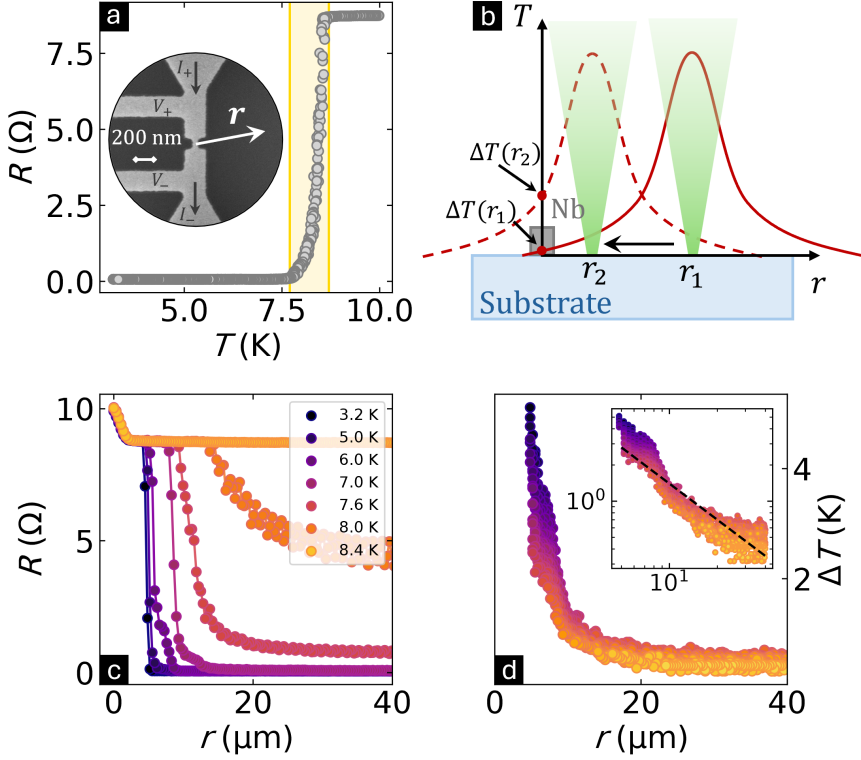


Figure 4.3: (a) Resistance of the junction as a function of temperature. The range over which the resistance-to-temperature conversion is exploitable is shaded in yellow. The inset shows a scanning electron microscope image of the probing device: the current is injected through the top and bottom terminals and the voltage is measured across the left-hand lines, and the arrow r defines the reference used in panels (c) and (d). (b) Schematic of the temperature profile induced by the absorbed laser light; the Nb junction measures the temperature rise as the laser is scanned across the surface. (c) Resistance as a function of the distance between the center of the laser spot and the center of the junction, for various base temperatures. The rise occurs when the laser reaches a distance at which the temperature profile drives the junction through its superconducting-to-normal transition at T_c ; the additional bump at distances of a few microns corresponds to the laser directly illuminating the junction and heating it far above T_c . (d) Reconstruction of the thermal profile, with each curve in (c) contributing one region of the profile, as indicated by the color coding. The inset shows the same data on a log-log scale, where the dashed line corresponds to $\Delta T \propto r^{-1}$.

a non-zero interface thermal resistance. This effect is balanced by the fact that the Nb junction reflects more light than the substrate which reduces thermal absorption.

In summary, despite the lack of resolution at short range, we are confident that the laser-induced heating is intense locally and extends far from the device at a weaker magnitude. This far-range characteristic distance is substrate-dependent through the thermal transport characteristics of the substrate.

Time dependence of the thermal profile

Because heat propagation is diffusive, the longer the heat is allowed to diffuse from an immobile source, the farther a significant temperature increase is observed. This behavior is governed by the heat diffusion equation

$$C_v \rho \frac{\partial T}{\partial t} = \kappa \frac{\partial^2 T}{\partial x^2} + \dot{q}, \quad (4.3)$$

where $\dot{q}$ is the quantity of heat introduced in the system per unit volume and time. From this equation, a characteristic time can be defined,

$$\tau \propto \frac{l^2}{D}, \quad (4.4)$$

with D the thermal diffusivity of Eq. (4.2) and l a given characteristic distance. This scaling, the steady-state $1/r$ profile measured in Fig. 4.3(d), and the diffusion length that governs the chopped-laser response of the following Section all follow from the same equation, as shown in Box 4.1. This characteristic time can be interpreted in two ways:

- It is the time required to heat up significantly an area of size l^2 . In the particular case of laser heating, a laser spot of finite size acts as the source, and the optical spot reaches its saturation temperature after the laser has been turned on for a time proportional to the area of the spot. In practice, given the

micron-range size of the spot and for the substrates used in this thesis, these timescales range from nanoseconds to microseconds.

- If the laser is chopped before entering the galvo scanner, the temperature of the sample oscillates at that same frequency. Far from the spot, however, this modulation is so strongly damped that it becomes negligible (as discussed in the following paragraph). When chopping at a frequency $f = 1/\tau$, the area over which the temperature effectively oscillates is given by $l \propto \sqrt{D\tau}$. Therefore, the faster the chopping, the smaller the perturbed area. This effect has been used extensively to improve the spatial resolution of numerous SLM studies [92, 98, 99].

Box 4.1 – Characteristic length and time in heat diffusion

Neglecting the source term, the heat diffusion equation reads $\partial_t T = D \nabla^2 T$, with $D = \kappa/\rho C_v$ (Eq. (4.2)). The Laplacian carries dimensions of $(\text{length})^{-2}$ and ∂_t of $(\text{time})^{-1}$, so any solution must relate a length l and a time t through the single scaling

$$l^2 \sim Dt, \quad (4.5)$$

which underlies the whole thermal behavior of the technique. Read as a time, it gives the characteristic timescale $\tau \sim l^2/D$ of Eq. (4.4); the two results below read it, instead, as a length.

Diffusion length. For a source modulated at angular frequency ω , the natural time is $t \sim 1/\omega$, and Eq. (4.5) reads as a length,

$$\lambda = \sqrt{\frac{2D}{\omega}}, \quad (4.6)$$

the *thermal diffusion length*: the distance over which a temperature oscillation is attenuated by $1/e$ [97]. It sets the damping and phase factors $e^{-r/\lambda}$ appearing in the chopped-laser response of the following Section.

Steady point source. The measured $1/r$ profile (Fig. 4.3(d)) follows from the same equation in the static limit. In three dimensions, a steady point source releasing power $\dot{Q}$ produces a radial profile obeying $\nabla^2 T = 0$ away from the origin; with $\nabla^2 T = r^{-2} \partial_r(r^2 \partial_r T)$, the bounded solution vanishing at infinity is

$$\Delta T(r) = \frac{\dot{Q}}{4\pi\kappa r}, \quad (4.7)$$

obtained by matching the total flux $-\kappa \oint \nabla T \cdot d\mathbf{S} = \dot{Q}$ through a sphere of radius r . The prefactor is set by κ alone: the local temperature rise reflects the thermal conductivity, whereas the diffusivity D controls only how the profile spreads in space and time.

Thermal Waves Induced by Chopped Laser Heating

To examine in greater detail what happens when a sample is locally heated by a chopped laser, we must turn to the framework of thermal waves. Even though heat is a diffusive process and therefore does not propagate like a wave, it can be shown that the *modulation* of the temperature at a given location does propagate as a wave. It is nonetheless not a wave in the strict sense, as it does not carry energy [100].

Thermal waves are generally analyzed in the harmonic regime, in which heat is delivered to the sample sinusoidally. In this case, the temperature distribution at the surface of a semi-infinite substrate ($w \gg d$, where w is the width and d is the thickness) is given by

$$\Delta T(r, t) = \text{Re} \left[\frac{I_0}{4\pi r \kappa} e^{-r/\lambda} e^{i(\omega t - r/\lambda)} \right], \quad (4.8)$$

where $\Delta T(r, t)$ is the temperature rise at radial position r from the source and time t , I_0 is the amplitude of the power input, ω is the

angular frequency of the excitation, and $\lambda = \sqrt{2D/\omega}$ is the thermal diffusion length (see Box 4.1, Eq. (4.6)) [97].

When a mechanical chopper is employed, the heat input into the system is approximated by a square wave that modulates the power between 0 and the laser power absorbed by the sample, I_0 . This input can therefore be expanded as a Fourier series of harmonic components:

$$I_{\text{abs}}(t) = \frac{I_0}{2} + \frac{2I_0}{\pi} \sum_{n=1,3,5,\dots}^{\infty} \frac{1}{n} \sin(n\omega t). \quad (4.9)$$

By substituting this modulation into Eq. (4.8), the resulting temperature response is expressed as

$$\Delta T(r, t) = \frac{I_0}{4\pi r \kappa} \left[\frac{1}{2} + \frac{2}{\pi} \sum_{n=1,3,5,\dots}^{\infty} \frac{1}{n} e^{-r\sqrt{\frac{n\omega}{2D}}} \sin\left(n\omega t - r\sqrt{\frac{n\omega}{2D}}\right) \right]. \quad (4.10)$$

This expression corresponds to a low-pass-filtered representation of a square wave, in which the filter exponentially attenuates the higher-order harmonics over distance. Several distinct dynamic regimes emerge from this relation, which are illustrated in Fig. 4.4. The top row displays the normalized temperature modulation ΔT as a function of the reduced time t/τ within one period of the excitation, together with the ideal square-wave drive delivered by the chopper (green dashed line). The bottom row shows the corresponding amplitude of the Fourier coefficients $|C_n|$ as a function of the harmonic number n , compared with the $1/n$ decay characteristic of an ideal square wave.

For $r \ll \lambda = \sqrt{\frac{2D}{\omega}}$, the exponential term in the summation approaches unity and the phase shift of the sinusoidal function is negligible for the dominant harmonics. Under these conditions, the summation reproduces the Fourier series of a square wave, so that the local temperature profile closely tracks the square profile of the excitation source. This is apparent in Fig. 4.4(a), where the temperature

follows the drive almost instantaneously, and in Fig. 4.4(d), whose Fourier amplitudes lie along the $1/n$ line up to high harmonic orders.

At intermediate distances, where $r \approx \lambda$, the Fourier coefficients are attenuated by a factor $e^{-\sqrt{n}}$. At the same time, the higher harmonics acquire distinct phase shifts $-r\sqrt{\frac{n\omega}{2D}} \approx -\sqrt{n}$ relative to the lower-order ones, misaligning their respective maxima. The superposition of these two effects introduces a pronounced time asymmetry in the temporal response, yielding a saw-tooth waveform characterized by a rapid relaxation immediately after a transition of the source state, followed by a gradual stabilization. This behavior is shown in Fig. 4.4(b): the response now lags the drive by a slight phase delay and loses its abrupt transitions, while the Fourier amplitudes in Fig. 4.4(e) fall increasingly below the $1/n$ line as the exponential damping suppresses the higher harmonics.

Finally, in the asymptotic limit $r \gg \lambda$, the high-order harmonics are strongly damped, leaving only the fundamental component; in this regime, the temperature undergoes low-amplitude, purely sinusoidal oscillations. As illustrated in Fig. 4.4(c), the surviving fundamental is subject to a large phase delay relative to the excitation, so that its maximum is markedly shifted in time with respect to the rising edge of the drive. Consistently, only the first few harmonics remain measurable in Fig. 4.4(f), the rest having been extinguished by the exponential attenuation.

Notably, these anharmonic thermal waves offer, in theory, the distinct advantage of decoupling the respective contributions of κ and D . In the near-field regime, the thermal diffusivity has a negligible influence on the local maximum temperature, which is predominantly governed by the prefactor inversely proportional to κ . This happens because the maximum temperature at the optical spot position results from a competition between heat generation (the laser power) and heat extraction (the thermal conductivity). Conversely, the diffusivity, together with the modulation frequency, sets the characteristic damping length of the oscillations. This spatial decay therefore determines the distance over which the thermal response transitions from a square

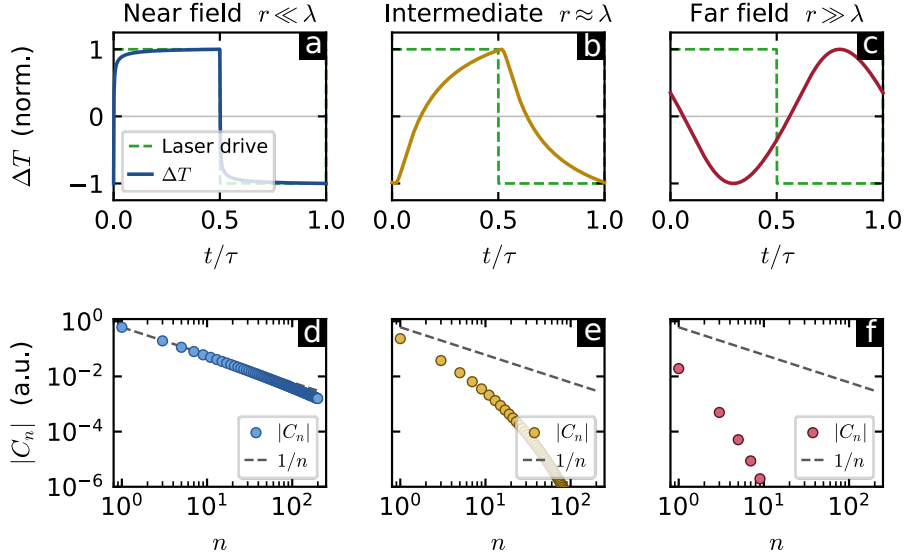


Figure 4.4: Temperature modulation induced by a chopped laser as a function of the reduced distance r/λ , computed from the Fourier expansion of Eq. (4.8). (a-c) Temperature response ΔT over one period of the excitation (solid line), together with the ideal square-wave laser drive (green dashed line). (d-f) Amplitude of the associated Fourier coefficients $|C_n|$ versus harmonic number n , compared with the $1/n$ decay of an ideal square wave (dashed line). (a,d) In the near-field regime ($r \ll \lambda$), the temperature closely retraces the square-wave drive and the coefficients follow the $1/n$ law up to high orders. (b,e) At intermediate distances ($r \approx \lambda$), the response acquires a slight phase delay relative to the drive and develops a saw-tooth shape, with a rapid relaxation after each transition followed by a gradual stabilization; the higher harmonics are progressively damped below the $1/n$ line. (c,f) In the far-field regime ($r \gg \lambda$), only the fundamental survives, giving a low-amplitude sinusoidal oscillation subject to a pronounced phase delay with respect to the excitation, and all higher harmonics are strongly suppressed.

wave to a saw-tooth profile, and asymptotically to a purely sinusoidal waveform.

4.3 Periodic time reconstruction via multiple harmonics collection

The multi-harmonic reconstruction developed here relies on a lock-in amplifier (LIA), which extracts the complex Fourier coefficient C_n of a signal at the n -th harmonic of a reference frequency f ; its operating principle is recalled in Box 4.2. For the particular case of the Stanford Research SR830 LIA used in this work, the harmonic index n is restricted to values satisfying both $n < 19999$ and $nf < 102$ kHz. It is worth noting that the technique is restricted to perfectly periodic $\text{Pr}(x, y)$, and that the time-resolved reconstruction does not work for samples exhibiting hysteretic behavior.

Box 4.2 – Lock-in amplifier and Fourier-coefficient extraction

A lock-in amplifier contains an internal oscillator that generates two reference signals phase-shifted by 90° , $R_{x,n}(t)$ and $R_{y,n}(t)$, which can be phase-locked to an external TTL trigger signal. Assuming that no phase shift relative to the sample excitation is configured and that the trigger event from the TTL signal at $t = t_0$,

$$R_{x,n}(t) = \sin(2\pi n f(t - t_0)) \quad (4.11)$$

$$R_{y,n}(t) = -\cos(2\pi n f(t - t_0)), \quad (4.12)$$

where f is the frequency of the reference signal and n is an integer multiple of the fundamental frequency f (*i.e.* the n -th harmonic). The measured voltage signal $V_{in}(t)$, which can be single-ended or differential and DC- or AC-coupled, is multiplied by each reference signal and passed through a low-pass filter of user-defined time constant τ_{LIA} . This analog signal processing is

the mathematical equivalent of

$$V_{x,n}(\mathcal{T}) = \int_{\mathcal{T}-\tau_{LIA}}^{\mathcal{T}} V_{in}(t) R_{x,n}(t) dt, \quad (4.13)$$

and similarly for V_y , using $R_{y,n}$. For $\tau_{LIA} \gg 1/(nf)$, the output of the LIA corresponds to the n^{th} Fourier coefficients of the input signal. Since the LIA uses a sine reference to measure the in-phase signal, the complex Fourier coefficients needed to compute the inverse transform are given by

$$C_n = -V_{y,n} + iV_{x,n}, \quad (4.14)$$

where C_n is the Fourier coefficient of the n^{th} harmonic and $i^2 = -1$.

Collecting a large number of harmonics allows the inverse Fourier transform to be computed accurately pixel-by-pixel. The result is the time evolution of the local thermal susceptibility of the SUT within one period of the excitation, $\tau = 1/f$. For N collected harmonics, $2N$ frames spaced in time by $\delta t = (2Nf)^{-1} = \tau/2N$ are obtained. Because the reconstructed signal must be purely real, the imposed Hermiticity adds a factor 2, effectively doubling the temporal resolution. Therefore, irrespective of the chopping frequency, the maximum achievable time resolution is set by the highest harmonic the LIA can measure, which in our case corresponds to $\delta t_{min} \approx 1/(2 \times 102 \text{ kHz}) \approx 4.9 \mu\text{s}$. In principle, a similar procedure could be carried out with an oscilloscope, but with much poorer voltage resolution.

As previously emphasized, this time-domain 2D reconstruction should not be misinterpreted as a direct visualization of the local thermal relaxation that follows localized heating. Although the reconstructed frames are evenly spaced in time, the mapped quantity is the macroscopic dynamic susceptibility $\text{Pr}(x, y)$ to the applied local stimulus. The evolution recorded at a given pixel is therefore a proxy for the rate of thermal diffusion toward the most sensitive regions

of the sample, parameterized by the phase shift relative to the periodic excitation. In essence, it is a transformation from the frequency domain rather than a strict spatiotemporal thermograph.

To validate the technique, we first apply it to reconstruct the time evolution of light reflected from a mirror. In this case, the detector is a photodiode, which is expected to record a square-wave photoresponse resulting from a laser source modulated by a mechanical chopper.

The time-domain voltage across the reverse-biased photodiode, measured with the oscilloscope, is shown in Fig. 4.5 for a chopping frequency $f = 188.27$ Hz. To obtain the frequency-domain response, the photodiode output is then connected to the LIA input, and the values of $V_{x,n}$ and $V_{y,n}$ are collected for all accessible harmonics. The sensitivity of the instrument is adjusted progressively as the amplitudes of the Fourier coefficients decrease. The resulting complex Fourier coefficients, plotted in a polar graph as amplitude and phase, are shown in Fig. 4.5(b). Because the inverse FFT takes as input the coefficients in both the negative and positive frequency domains, the complex Fourier coefficients are extended in a Hermitian manner, ensuring that the temporal reconstruction is a purely real function. The DC coefficient at $f = 0$ Hz cannot be measured by the LIA and must instead be obtained with a voltmeter to fully reconstruct the FFT. Its effect, however, is merely to shift the IFFT vertically, so that it carries no information about the temporal dynamics of the response. It is measured in this example but subsequently ignored in the case study presented in the second part of Chapter 6. Fig. 4.5(c) shows the amplitude of the Fourier coefficients. As expected for a 1/2 duty-cycle square wave, the amplitude decreases inversely proportionally to the odd harmonic number, whereas the even harmonics are negligibly small. Fig. 4.5(d) shows the reconstructed signal (red dots), obtained from the IFFT of the frequency-domain measurements, plotted on top of the time-domain signal of panel (a). The inset in panel (c) shows a close-up of the downward transition in the photodiode voltage, emphasizing the excellent fidelity of the reconstruction to the original signal.

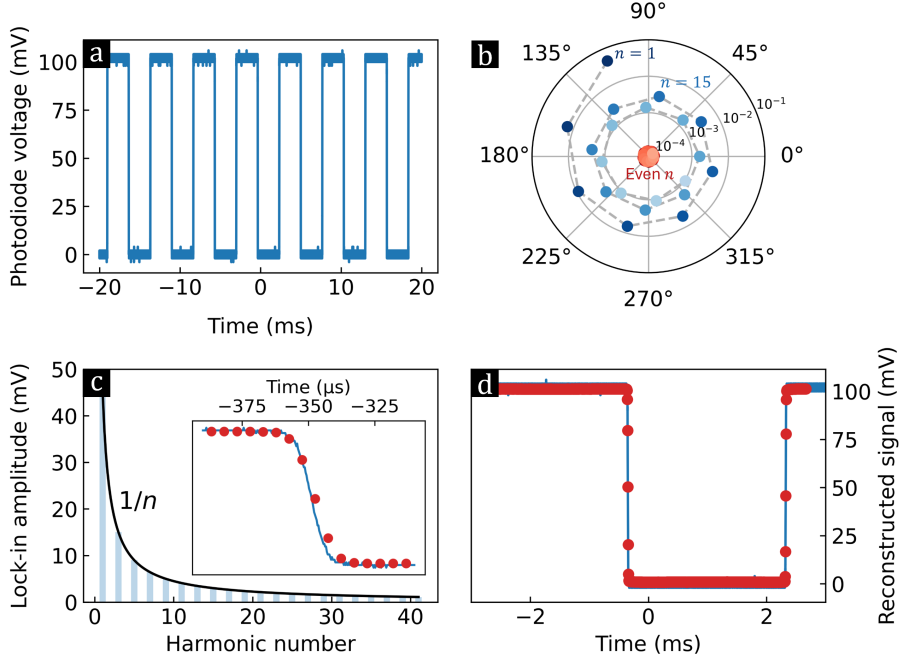


Figure 4.5: (a) Example signal measured with an oscilloscope. The laser is chopped at a fixed frequency $f = 188.27$ Hz and the reflected light is collected by a photodiode. The signal is periodic and can therefore be reconstructed by the inverse Fourier transform method. (b) Fourier coefficients measured with the LIA for the first 40 harmonics, shown in the complex plane. The fundamental ($n = 1$) is the strongest, the amplitudes of the odd harmonics decrease inversely with harmonic number, and the even harmonics are markedly weaker; the phase encodes the synchronization of the detected signal with the reference TTL signal. (c) Amplitude of the measured coefficients. (d) Reconstructed signal (red) overlaid on the signal of panel (a). The inset in panel (c) shows a close-up of the downward transition in the photodiode voltage, highlighting the excellent fidelity of the reconstruction to the original signal.

Chapter 5

Magnetic Flux Dynamics in Low- T_c Superconducting Resonators

Planar superconducting resonators operating at radio-frequencies have driven major advances in a wide range of applications, including quantum sensing and detection [1, 5], quantum computing [2, 3], and information processing [4, 6]. Their unique advantage stems from the zero electrical resistance of the superconducting materials composing these resonators. The aforementioned applications require low energy losses, characterized by high quality factors Q , and superconducting resonators have been shown to reach quality factors up to nearly 10^7 , several orders of magnitude higher than their metallic counterparts [1].

To achieve such performance, it is crucial that the superconducting resonator remains free of superconducting vortices, as their motion leads to significant energy dissipation and deteriorates the quality factor [8]. Although magnetic screening methods can be used to mitigate this problem, some applications require exposure to an external magnetic field [101]. Some qubit implementation schemes such as spin ensembles in solid-state systems [13–15], phase-slip qubits [16–18, 102], and trapped electrons [19, 20] require inevitable exposure to a magnetic

field. In addition, even though planar superconducting resonators are resilient to in-plane magnetic fields [21], unwanted small in-plane misalignments can give rise to out-of-plane magnetic field components that can severely affect their performance. For all these reasons, the identification of new strategies and designs for increasing the magnetic field resilience of superconducting CPW resonators remains a subject of major technological significance.

The deleterious effect of the magnetic field can be attributed to the penetration of magnetic flux quanta, or vortices [7]. Their interaction with the induced radio-frequency (RF) currents gives rise to energy absorption, resulting in dissipation that translates into a decrease in the quality factor and the resonance frequency [8–11, 103, 104]. Not only the total number of vortices penetrating the CPW resonator is relevant, but also their exact location [105]. Two broad approaches can be adopted to lessen vortex-induced losses: (i) hinder their movement by creating artificial pinning sites [9, 11, 22, 103, 104], and (ii) decrease their number in the critical regions by, for example, reducing the ground plane [23, 106] or dividing it into fractal structures [24, 25].

At temperatures below approximately half the superconducting critical temperature T_c , these strategies are further complicated by the emergence of thermomagnetic instabilities. These instabilities can trigger magnetic flux avalanches, the rapid nucleation and propagation of magnetic flux into the superconducting material [29], which, as this chapter will show, introduce stochastic, unpredictable behavior manifesting as fluctuations in the resonance frequency as a function of the applied magnetic field. Strategies based on increasing vortex pinning, such as through holes, may actually favor the triggering of magnetic flux avalanches at low temperatures and thereby deteriorate the performance [35, 36].

In order to seize and quantify the impact of these various phenomena, the complete sample surface must be taken into account. Magneto-optical imaging (MOI) has proven to be an excellent tool for examining macroscopic structures such as superconducting CPW resonators with μm spatial resolution [30, 38]. A one-to-one com-

parison between MOI observations and high-frequency transmission measurements allows direct identification of the non-local magnetic origin of the features observed in $f_r(\mu_0 H_a)$.

An important open question is whether the RF excitation itself stimulates the nucleation and propagation of magnetic flux avalanches. From a theoretical standpoint, at temperatures close to the critical one, the first-vortex penetration field for RF magnetic fields is somewhat larger than for DC fields [107, 108], which could be beneficial for keeping vortex penetration at bay. However, this difference is expected to become less pronounced at low temperatures [109, 110]. Experimentally, some studies have suggested that RF excitation promotes flux avalanches [30–32], while others report suppression of avalanche activity for probing frequencies near the depinning frequency [33]. Numerical simulations show that AC excitation reduces the threshold field for avalanche onset [34]. The literature thus remains inconclusive, partly due to the lack of direct experimental evidence.

In this chapter, we first investigate the effect of magnetic field penetration on the resonance frequency of NbTiN superconducting resonators by combining *sequential* MOI and high-frequency measurements (Sections 5.4 and 5.5). A clear and unambiguous correspondence between both independent methods is demonstrated. We then address the question of RF-induced avalanches (Section 5.6) by upgrading the setup to perform *simultaneous* RF transmission and MOI measurements, enabling for the first time the unambiguous tagging of individual avalanche events to specific jumps in resonance frequency. Electromagnetic simulations further clarify the influence of avalanche location on the resonance response.

A note on the structure of this chapter. The results presented here were obtained over two experimental investigations, separated in time. The first, adapted from L. Nulens & N. Lejeune *et al.*, “Catastrophic magnetic flux avalanches in NbTiN superconducting resonators” [28], relied on *sequential* measurements: the magnetic landscape and the RF transmission were acquired in separate cooldowns and setups

on twin samples (Device #1). For the second study, adapted from J. Baumgarten & N. Lejeune *et al.*, “Magnetic landscape of NbTiN superconducting resonators under radio-frequency excitation” [48], the experimental setup was upgraded to allow *simultaneous* magneto-optical imaging and RF transmission on a single device (Device #2), which required removing the metallic mirror of the magneto-optical indicator to avoid parasitic eddy currents. The chapter follows this chronology: Sections 5.3, 5.4 and 5.5 present the sequential study, and Section 5.6 the simultaneous one. Because both works address the same physical system, some concepts (device layout, temperature characterization, the four flux-penetration regimes) recur across the two parts; they are introduced once, in the sequential part, and only the genuinely new elements are developed in the simultaneous part.

Personal contribution. In the first study [28], the samples were fabricated at Chalmers University of Technology and the radio-frequency characterization was carried out at KU Leuven, where it was later reproduced in Liège after publication. I carried out the magneto-optical imaging and its analysis. The correspondence between the magnetic landscape and the field-dependent resonance frequency, and its interpretation, were established collectively with the co-authors. For the second study [48], I adapted the cryostat and the magneto-optical setup to make them compatible with radio-frequency measurements, which included the removal of the metallic mirror of the indicator discussed in Section 5.6.2. The measurement campaign was conducted jointly with J. Baumgarten, who took the lead on the data analysis; the interpretation and the discussion were a collective effort, and the respective contributions of both authors are essentially equivalent.

5.1 Sample design

Two device geometries, shown in Fig. 5.1, are used throughout this chapter. **Device #1** (Fig. 5.1(a)) is used for the sequential MOI

and RF measurements described in Sections 5.4 and 5.5. Several identical samples were fabricated to allow subsequent comparison between MOI and RF measurements. In order to fit the complete resonator structure within the field of view of the MOI setup, a substrate size of $(5 \times 4) \text{ mm}^2$ was chosen, with a superconducting region of $(2.5 \times 2.5) \text{ mm}^2$. Two overcoupled, hanger-type $\lambda/4$ resonators with lengths $4089 \mu\text{m}$ (shorter) and $3953 \mu\text{m}$ (longer) are capacitively coupled to a central feedline. The central conductor has a width $w = 20 \mu\text{m}$ with a separation from the ground plane of $g = 10 \mu\text{m}$, resulting in a 50Ω impedance. The resonators are meander-like, with five U-turns.

Device #2 (Fig. 5.1(b)) is used for the simultaneous RF-MOI measurements described in Section 5.6. It contains three overcoupled $\lambda/4$ resonators (labeled #1, #2, and #3) with lengths $4392 \mu\text{m}$, $4561 \mu\text{m}$, and $4744 \mu\text{m}$, capacitively coupled to a central feedline with $w = 20 \mu\text{m}$ and gap $g = 10 \mu\text{m}$. In contrast to Device #1, the resonators in Device #2 are limited to a single U-turn and each includes an unused DC flux-biasing line.

5.2 Experimental methods

5.2.1 Sample fabrication

The resonators of Device #1 were fabricated on double-side polished C-plane Al_2O_3 wafers of thickness $(430 \pm 25) \mu\text{m}$, or on $525 \mu\text{m}$ -thick highly resistive Si for Device #2. The wafer was first cleaned using a standard 5:1:1 RCA process of deionized water, ammonia water, and H_2O_2 heated to 80°C for 10 min [111]. A 100 nm NbTiN film was then deposited by sputtering from a $\text{Nb}_{81.9}\text{Ti}_{18.1}\text{N}$ target [112] in a near-UHV sputter system under an $\text{Ar}:\text{N}_2$ partial gas pressure ratio of 60:37, with in-situ annealing at 660°C for 30 min prior to and during deposition. The coplanar waveguide resonators were patterned using electron-beam lithography (PMMA A4 resist with an E-spacer conducting layer to mitigate charging effects), followed by Ar/Cl

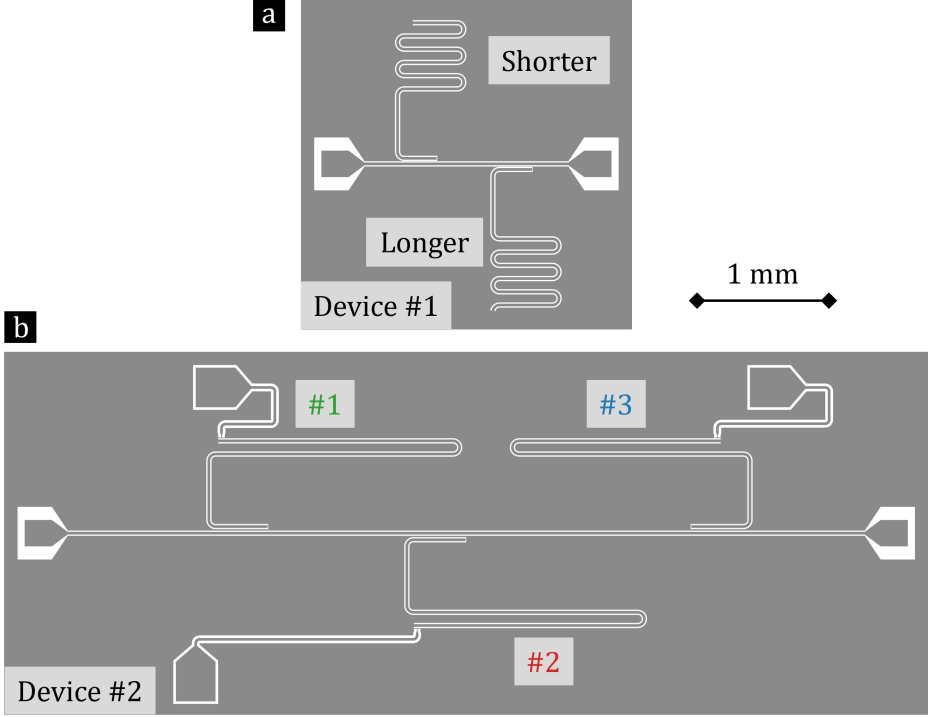


Figure 5.1: Layout of the two devices studied in this chapter. (a) Device #1, with two meander-like $\lambda/4$ resonators (shorter and longer) of five U-turns each, capacitively coupled to a central feedline. (b) Device #2, with three single-U-turn resonators (#1, #2, and #3) coupled to a central feedline, each terminated by an unused DC flux-biasing line; the label colors match the color code used for the resonators throughout the chapter. The superconducting NbTiN is shown in gray and the etched regions in white; the scale bar is common to both panels.

(4:50 sccm) ICP plasma etching [3, 113]. Residual resist was removed by heated acetone. A critical temperature of 16.5 K was obtained using SQUID magnetometry on a reference layer of 100 nm. The fabrication was carried out at MyFab, Chalmers University of Technology.

5.2.2 RF measurements

High-frequency measurements were performed using a Rohde & Schwarz ZNL20 Vector Network Analyzer (VNA) for Device #1 and a Keysight P5003B Streamline VNA for the second. For Device #1, a Janis ^3He cryostat was used; the input line was attenuated with 60 dB distributed across different temperature stages. The sample was mounted on an oxygen-free copper holder and connected via wire bonding to an RF holder with a sample space of $(7 \times 4) \text{ mm}^2$. A 4–8 GHz circulator/isolator and 4–8 GHz high-electron-mobility transistor (LNF-LNC4_8C) were connected at the 4 K stage after the output line of the sample. The magnetic field was applied perpendicular to the sample using a superconducting magnet with a sweep rate of $\sim 184 \text{ Oe/s}$. To compensate for the remanent field, the sample was field-cooled (-0.2 mT) before every measurement cycle. A measurement cycle consists of sweeping the magnetic field along the sequence $0 \rightarrow 5 \rightarrow -5 \rightarrow 5 \rightarrow 0 \text{ mT}$. Initial field cycles were executed at -40 dBm driving power and compared to 0 dBm ; since no significant difference was observed, all measurements were executed at 0 dBm . For the simultaneous RF-MOI measurements with Device #2, a Montana Instrucent closed-cycle cryostat *Cryostation C1* was used and the magnetic field was applied via a copper coil with a sweep rate of $\sim 44 \text{ mT/s}$. The VNA acquires the forward transmission coefficient S_{21} ; the resonance frequency f_r and quality factors are extracted via the fitting routine described in Chapter 2.

5.2.3 Magneto-optical imaging

The magnetic flux landscape was visualized by magneto-optical imaging (MOI), whose physical principle, setup, and post-processing are described in detail in Chapter 3. In short, a 3 μm -thick Bi-doped yttrium iron garnet (Bi:YIG) indicator film is placed in direct contact on top of the sample and the Faraday rotation of linearly polarized light yields images whose intensity is proportional to the local out-of-plane field B_z . Crucially for the simultaneous measurements of Section 5.6, the standard indicator carries an aluminum mirror on its back. This metallic layer is removed in the upgraded configuration, as discussed in Section 5.6.2.

5.2.4 SQUID magnetometry

Magnetic measurements were carried out in a SQUID magnetometer (Quantum Design SQUID VSM MPMS3). The DC magnetization versus temperature was measured in a parallel applied field upon heating from 1.8 to 300 K under a remanent magnetic field of the order of 0.01 mT. The critical temperature T_c was defined as the highest temperature with a measurable diamagnetic magnetization contribution. A demagnetization procedure was conducted at 300 K prior to each measurement by oscillating the field from -7 T to near zero to minimize the remanent field from the superconducting magnet [114].

5.3 Temperature-dependent resonator characterization

The resonance frequencies f_r of the two resonators in Device #1, at 7.06 GHz and 7.44 GHz, manifest themselves as absorption dips in the S_{21} transmission spectrum, with internal quality factors of ~ 2000 and ~ 800 extracted from the fitting routine described in Section 2.1.3 [61]. The dips show a pronounced skewness caused by the impedance

mismatch caused by the lengthy wire bonds ($\sim \text{mm}$) required by the high-frequency setup [60, 115, 116]. As the temperature increases, f_r shifts to lower values and the linewidth broadens, indicating increasing losses.

As recalled in Section 2.1.4, the resonance frequency of a superconducting $\lambda/4$ resonator (Eq. (2.23)) carries, beyond its geometric contribution, a temperature-dependent kinetic-inductance term $L'_k(T)$ that diverges at T_c (Eq. (2.22) and Box 2.6). Fitting the measured relative frequency shift $\delta f = [f(0) - f_r]/f(0)$ with L'_{k0} , T_c , and L'_g as free parameters yields the values gathered in Table 5.1.

Resonator	L'_g (nH/mm)	L'_{k0} (nH/mm)	T_c (K)
Shorter	0.382 ± 0.0005	0.0373 ± 0.00006	17.27 ± 0.02
Longer	0.389 ± 0.0007	0.0511 ± 0.00009	17.21 ± 0.02

Table 5.1: Fitting parameters of the two resonators of Device #1, obtained from the temperature-dependent resonance frequency (Eqs. (2.23) and (2.22)).

The geometric inductance is relatively close to its design value $L'_g = 0.419 \text{ nH/mm}$, and the excellent fit (not shown for Device #1) confirms that the resonance frequency serves as a sensitive probe for any variation of the superconducting energy gap.

5.4 Magnetic field response at low temperature

5.4.1 Different regimes of flux penetration

Let us now investigate the response of the resonators as a function of the applied out-of-plane magnetic field $\mu_0 H_a$. Fig. 5.2(a) shows the resonance frequency of the longer resonator of Device #1 after a zero-field cooling process and subsequent sweep $0 \rightarrow 5 \rightarrow -5 \rightarrow 0 \text{ mT}$ at a fixed temperature of 5 K. For clarity, the cycle has been unfolded so

that the abscissa can be associated with the measurement chronology. As the magnetic field increases, f_r shifts down at a rate smaller than the subsequent increase rate when decreasing field. This irreversible, hysteric behavior of f_r is a universal feature observed for Nb [104, 117], NbN [118], YBCO [119], MgB₂ [120], and Al [121] CPW resonators, and can be linked to the inhomogeneous distribution of RF current I_{rf} shaking an inhomogeneous magnetic flux distribution. As shown in Supplementary Note 1 of Ref. [28], the internal quality factor Q_i has a very similar magnetic field dependence, demonstrating that the same interplay defines dissipation in these resonators.

A close inspection of Fig. 5.2(a) reveals several distinct features, labeled with Roman numerals (I–IV), which are directly correlated to the MOI observations in panel (d).

Magnetic screening regime (I). In the field range $0 < \mu_0 H_a < 0.4$ mT, screening currents circulate in the ground plane. The ground plane and resonators are capacitively coupled to the excitation line and therefore physically disconnected from it, making the structure topologically equivalent to a ring surrounding a disk. Screening currents circulating clockwise in the ground plane produce a magnetic field of opposite polarity at the inner border of the ground plane and a small positive field inside the central hole (where the feedline sits). As long as $0 < \mu_0 H_a < 0.4$ mT, the resonators are protected by the screening currents and f_r gently decreases as the field increases.

Magnetic perforation (transition between I and II). At $\mu_0 H_a \approx 0.5$ mT, a magnetic perforation event occurs: a sudden breakdown of the screening currents in the ground plane with simultaneous injection of magnetic flux into the central hole [122–124]. This event marks the start of a steeper field dependence of f_r . The injection of the magnetic field becomes apparent in the differential MOI image at this transition (Fig. 5.2(e)), where the nucleation point of the magnetic perforation is indicated with a dotted circle.

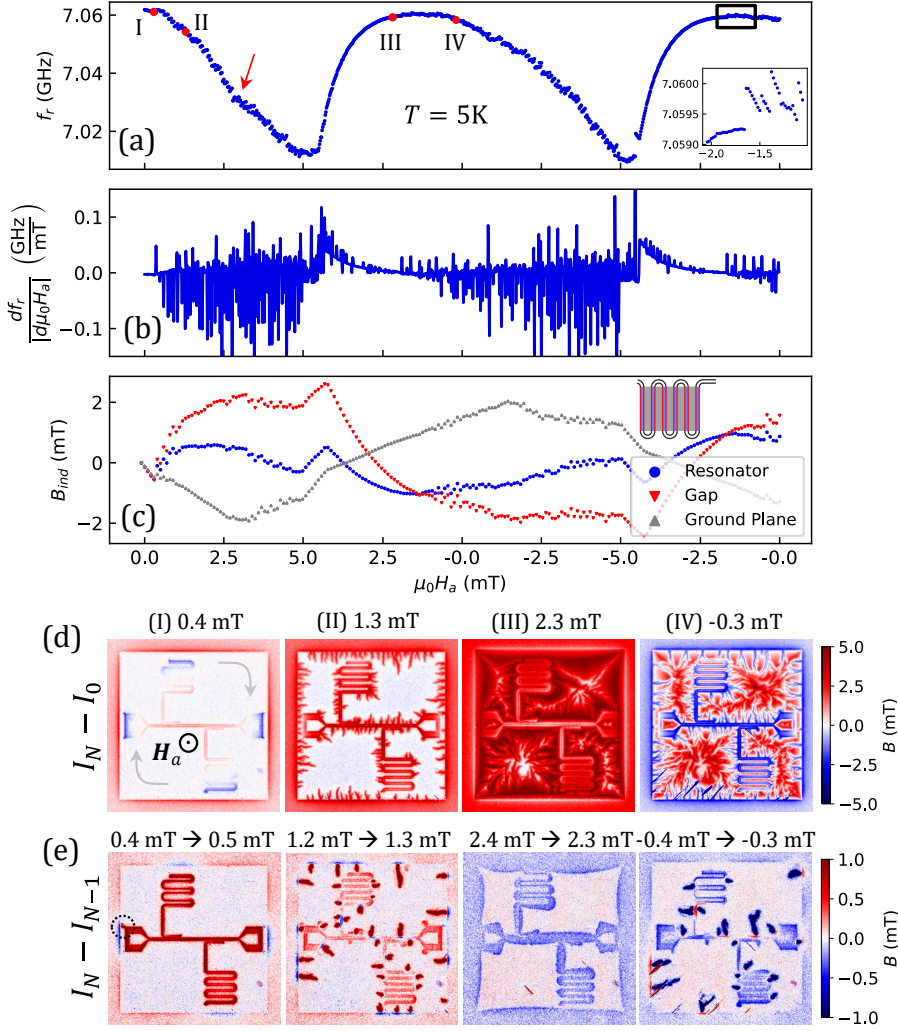


Figure 5.2: (a) Resonance frequency f_r along a full field loop $0 \rightarrow 5 \rightarrow -5 \rightarrow 0$ mT (inset: zoom of the framed region). Features I–IV are described in the text; the red arrow marks the kink. (b) Derivative $df_r/d\mu_0 H_a$, revealing the noisy and smooth regimes. (c) Induced field $B_{\text{ind}} = B - \mu_0 H_a$ averaged over the meander central conductor (blue), the gap (red), and the ground-plane pitch (gray), as sketched in the inset. (d) Magneto-optical images at the field values labeled I–IV in panel (a); gray arrows in image I indicate the magnetic screening currents. (e) Differential images between consecutive field steps; the dotted circle marks the magnetic perforation nucleation point.

Flux avalanche regime (II). As the field increases beyond the perforation point ($0.5 < \mu_0 H_a < 5 \text{ mT}$), $f_r(\mu_0 H_a)$ exhibits a noisy behavior with a global linear decrease, better seen in the derivative $df_r/d\mu_0 H_a$ shown in Fig. 5.2(b). The origin of this jumpy behavior is the development of magnetic flux avalanches caused by thermomagnetic instabilities [125]. A snapshot of the magnetic landscape showing finger-shaped flux avalanches at point (II) is shown in Fig. 5.2(d). These avalanches are neither fully reproducible [126, 127] nor totally random, and can originate at border spots with either high or low degrees of perfection [124, 128].

A kink in the middle zone of the noisy regime, indicated by a red arrow in Fig. 5.2(a), is also observed. To identify its origin, the local induced field $B_{\text{ind}} = B - \mu_0 H_a$ is measured in three regions: the central conductor of the resonator meander (blue), the gap separating the resonator from the ground plane (red), and the ground plane pitch between the meander structure (gray). The kink arises precisely when the induced field in the ground plane starts to increase, due to the development of flux avalanches in the ground plane. These events relieve the magnetic pressure everywhere in the resonator and slow down the decrease of f_r .

Smooth flux removal (III). When decreasing the magnetic field from $\mu_0 H_a = 5 \text{ mT}$, the local field inverts polarity at the sample border before the applied field changes sign [37, 129, 130]. Antivortices nucleate at the border and annihilate with the trapped vortices. Since no flux avalanches occur in this regime, $f_r(\mu_0 H_a)$ does not exhibit any jumpy behavior. The efficient removal of flux from the border due to the pronounced decrease of the total magnetic field at the sample's borders results in a lower vortex density during the downsweep than the upsweep at the same applied field, giving rise to a steeper increase of f_r as $\mu_0 H_a$ decreases. Although previously created avalanche channels leave a frozen imprint on the magnetic landscape that may act as channels for the penetration of antivortices [131], the differential imaging demonstrates that no new avalanches develop in this regime.

Anti-avalanche regime (IV). This smooth regime ends at a threshold field $\mu_0 H^* \approx 2.5$ mT. Decreasing the applied field below this value triggers avalanches of flux with polarity opposite to the applied field, accompanied by a jumpy $f_r(\mu_0 H_a)$ response down to $\mu_0 H_a = -5$ mT. After each anti-avalanche event, the magnetic pressure at the border of the resonator is relaxed and f_r increases abruptly; as $\mu_0 H_a$ continues to decrease, the system rebuilds magnetic pressure and f_r slowly decreases again, producing a characteristic sawtooth pattern visible in the inset of Fig. 5.2(a).

5.5 Temperature dependence of the avalanche activity

Flux avalanches have a critical impact on the overall performance of the CPW resonators. Fig. 5.3(a) shows the magnetic flux penetration after zero-field cooling for three temperatures (5, 9, and 12 K) at the same applied field $\mu_0 H_a = 1.1$ mT. Finger-like avalanches clearly visible at 5 K are no longer present at 9 and 12 K. The occurrence of thermomagnetic instabilities is known to be restricted to low temperatures (roughly $T < T_c/2$). At $T = 9$ K, border defects act as magnetic faucets from which magnetic flux is injected into the ground plane [128]. At $T = 12$ K, the sample is largely penetrated by the magnetic field and non-penetrated regions shrink to form discontinuity lines where the supercurrents change direction [132].

The threshold field $\mu_0 H^*$ at which anti-avalanches are triggered can be identified as the onset of the jumpy $f_r(\mu_0 H_a)$ response (shadowed region in Fig. 5.3(b)) or directly from MOI measurements. Fig. 5.3(c) shows $\mu_0 H^*$ as a function of temperature from both independent methods, for the decreasing (∇) and increasing ($\triangle$) branches of the hysteresis loops. To ensure consistency, the experiment was repeated during a second cooldown at 5 K and 9 K. As T increases, $\mu_0 H^*$ progressively diminishes and no avalanche activity is observed above 10 K. Both methods are in excellent agreement, validating the

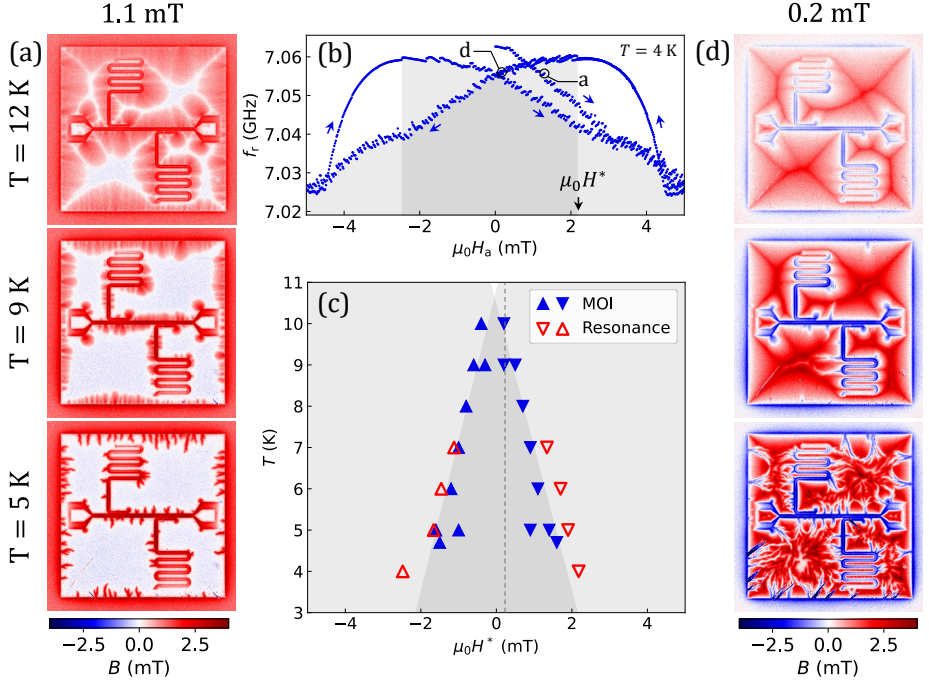


Figure 5.3: (a) Magnetic landscape after zero-field cooling at three different temperatures (5, 9, and 12 K) for the same applied field $\mu_0 H_a = 1.1$ mT. (b) Hysteresis loop of the resonance frequency $f_r(\mu_0 H_a)$ obtained at 4 K. The region where (anti-)flux avalanches are triggered when decreasing (increasing) from a maximum (minimum) applied field is highlighted with a gray background. (c) Threshold field $\mu_0 H^*$ as a function of temperature as determined by the onset of the noisy regime in $f_r(\mu_0 H_a)$ measurements (red triangles) and by MOI measurements (blue triangles), for the decreasing (∇) and increasing ($\triangle$) branches of the hysteresis loops. The dashed line indicates the field $\mu_0 H_a = 0.2$ mT at which the images in panel (d) were selected. (d) Magnetic landscape in the decreasing branch of the hysteresis loop at $\mu_0 H_a = 0.2$ mT for the same three temperatures, illustrating the presence of anti-flux avalanches at 5 and 9 K and their absence at 12 K.

use of RF transmission measurements as a non-invasive probe of thermomagnetic instabilities.

5.5.1 Strategies for magnetic field resilience

Based on the measurements and analysis presented above, we can devise a roadmap for design optimization of superconducting resonators with the goal of improving their resilience under perpendicular magnetic fields.

In order to increase the threshold field needed for magnetic perforation, it is essential to decrease the local field at the outer rim of the ground plane. In the magnetic screening state, this local field can be approximated as $\approx \sqrt{W/t} \mu_0 H_a$ [129], where W is the width of the ground plane and t is the film thickness. The desired effect can therefore be achieved by reducing the lateral dimension W of the ground plane [23, 106].

If flux avalanche-induced noise needs to be suppressed, a possible solution consists in reducing the pitch of the meanders. It has been shown that thermomagnetic instabilities are progressively suppressed as the width of the superconductor decreases and are fully absent below a certain threshold value [133], as exemplified in the field resilience of nanowire NbTiN resonators [134]. Alternatively, introducing artificial roughness along the borders of the superconducting device can shrink the H - T region where avalanches occur [128, 135]. The introduction of periodic pinning (through holes) can be beneficial for reducing losses caused by flux shaking, but only at sufficiently high temperatures where flux avalanches do not take place; at low temperatures, such pinning may actually favor avalanche triggering and deteriorate device performance [35, 36]. Another interesting approach consists in adding pinning centers in the resonator where the RF local field maximizes [41]. Finally, it may be worth exploring asymmetric pinning landscapes forming flux lenses [136], although the efficiency of such ratchet potentials demonstrated at relatively low frequencies seems less promising at GHz frequencies [137].

It is worth noting that the maximum applied magnetic field of 5 mT explored in this section may seem modest, but the required field intensities for various applications can greatly exceed this range. For instance, in solid-state spin ensemble systems the required field typically falls within 15–200 mT [13–15], for phase-slip qubits the range is around 0.3 mT [16, 17], whereas for trapped electrons it lies in the Tesla range [19, 20]. While the magnetic field is applied in-plane in some of these systems, a slight misalignment can introduce a significant perpendicular field component. It is therefore crucial to extend the present studies to higher fields in future work.

The sequential measurements presented so far establish an unambiguous correspondence between the magnetic flux landscape and the resonance frequency. However, they leave a central question open: does the RF excitation itself promote or stimulate magnetic flux avalanches? Answering this requires acquiring the magnetic landscape and the RF transmission *simultaneously*. As anticipated above, this endeavor is highly non-trivial, since the proximity of the metallic mirror and the magnetic indicator needed for MOI may substantially distort the performance of the resonators. In the following, we describe the setup upgrade that overcomes this difficulty and the new physics it reveals.

5.6 Simultaneous radio-frequency and magneto-optical measurements

The upgraded setup is sketched in Fig. 5.4(a): the three multiplexed $\lambda/4$ resonators of Device #2 are excited through the feedline of a VNA while, at the same time, the magnetic landscape is imaged in widefield through the Bi:YIG indicator placed on top of the sample. Beyond confirming the sequential picture, this configuration removes its main limitation being the impossibility of identifying which individual flux avalanche is the source of a given jump in resonance frequency. It allows us to address the causal relation between RF excitation and

avalanche activity in both directions.

5.6.1 Characterization of Device #2

The three resonances of Device #2, observed around 5.08, 5.42, and 5.72 GHz (Fig. 5.4(b)–(d)) and corresponding to resonators 1, 2, and 3, are fitted following Eq. (2.15). In contrast to Device #1, the presence of the DC bias lines distorts the resonance peaks, indicating additional losses and parasitic modes [138]. In particular, resonator 1 exhibits a marked asymmetric shape due to overcoupling, whereas resonator 2 shows a lower quality factor. Quality factors on the order of 10^3 – 10^4 are found for all bare resonators.

The relative frequency shift $\delta f_r = f_r(T) - f_r(T_0)$, with $T_0 = 3.33$ K, again follows from the temperature-dependent kinetic inductance (Eqs. (2.23) and (2.22)) and is fitted (Fig. 5.4(e)) with L'_g , C' , L'_{k0} , and T_c as free parameters. The resulting values for the three resonators are gathered in Table 5.2.

Resonator	L'_g (nH/mm)	C' (nF/m)	L'_{k0} (nH/mm)	T_c (K)
#1	0.544 ± 0.003	0.161 ± 0.008	0.1230 ± 0.0006	14.51 ± 0.45
#2	0.562 ± 0.003	0.152 ± 0.008	0.1080 ± 0.0005	14.61 ± 0.46
#3	0.421 ± 0.003	0.202 ± 0.002	0.0690 ± 0.0005	14.56 ± 0.62

Table 5.2: Fitting parameters of the three resonators of Device #2, obtained from the temperature-dependent resonance frequency (Eqs. (2.23) and (2.22)).

These T_c values are lower than those of Device #1, consistent with the samples having been fabricated at different times and a possible poorer thermal contact between the sample and the cold finger. They nonetheless confirm that the measurements performed at 3.2 K lie well within the avalanche regime ($T \lesssim T_c/2$).

The power dependence of the resonance dip (inset of Fig. 5.4(e)) reveals a gradual decrease in resonance frequency as the probing power increases. At RF powers exceeding -20 dBm, the transmission

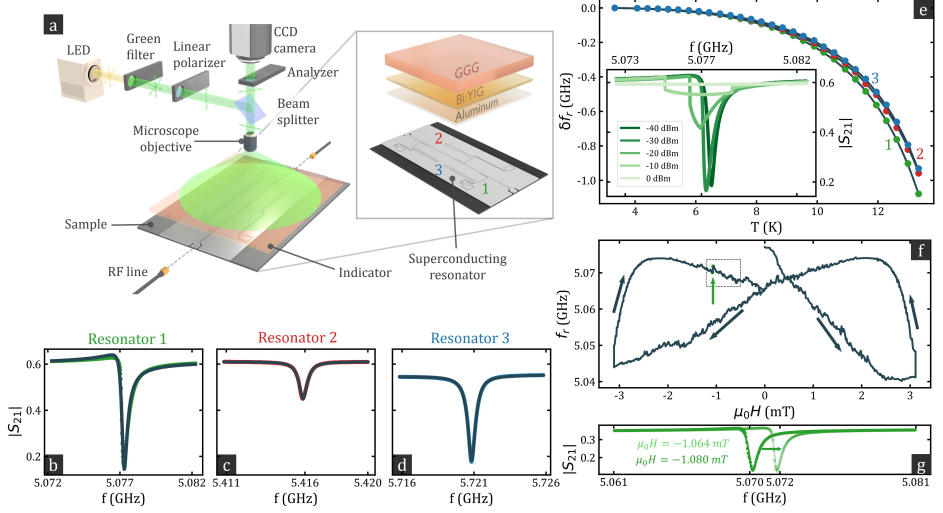


Figure 5.4: (a) Experimental setup combining RF excitation of three multiplexed $\lambda/4$ resonators and simultaneous magneto-optical imaging of the sample. (b)–(d) Transmission coefficient S_{21} in the vicinity of the three resonance frequencies at 3.3 K and a RF power of -30 dBm; black curves are fits. (e) Relative change of the resonance frequency, $\delta f_r = f_r(T) - f_r(T = 3.33 \text{ K})$, as a function of temperature; the lines are fits combining Eqs. (2.23) and (2.22). The inset shows the resonance peak shape as a function of excitation power. (f) Magnetic field dependence of the resonance frequency for an applied out-of-plane field cycled $0 \rightarrow 3.3 \rightarrow -3.3 \rightarrow 0$ mT, with dark arrows indicating the sweep direction. The green arrow marks the situation shown in panel (g); the dotted box indicates the field interval studied in Fig. 5.7. (g) Transmission parameter for two consecutive applied magnetic fields, evidencing a single avalanche-induced jump event.

spectrum shows a steep, discontinuous jump on the low-frequency side of the resonance dip, indicating entry into a nonlinear regime characterized by bifurcations in the S_{21} spectrum [139–141]; at 0 dBm the resonance peak is already entirely suppressed. This effect may arise from the power-dependent current and the nonlinear kinetic inductance. At elevated microwave powers (above 0 dBm), other mechanisms can develop, such as hot-spot formation [142, 143]. Increasing the RF power naturally leads to a larger RF magnetic field, which can eventually exceed the vortex penetration field (a few mT amplitude) and thereby produce microwave hysteresis losses, as discussed for $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ and NbN thin films in Ref. [144].

Finally, the resonance frequency is also susceptible to the out-of-plane component of the applied field. Sweeping the out-of-plane magnetic field while tracking the resonance frequency reveals the characteristic butterfly-like shape [104] shown in Fig. 5.4(f), in which $f_r(\mu_0 H_a)$ exhibits an irreversible response when cycling the field. The critical state model for hard superconductors naturally accounts for this behavior through the magnetic field dependence of the surface impedance and the effective RF field penetration [62, 145, 146]. A noisy response, in which the resonance frequency undergoes sudden jumps (one such jump is shown in Fig. 5.4(g)), is observed over an extensive range of applied fields. Such stochastic response was also reported for in-plane applied fields in Ref. [26], likely due to unwanted out-of-plane field components. It is worth mentioning that the actual magnetic field felt by the superconducting circuit can exceed the applied field by an order of magnitude or more due to the flux focusing effect [23, 147].

5.6.2 Invasiveness of the magneto-optical indicator

To perform MOI on the planar devices, an optically active indicator is placed on top of the sample. Unfortunately, the presence of the magnetic indicator significantly influences the transmission spectrum and the resonance parameters, as illustrated in Fig. 5.5.

In the standard MOI indicator configuration, the Bi:YIG film is

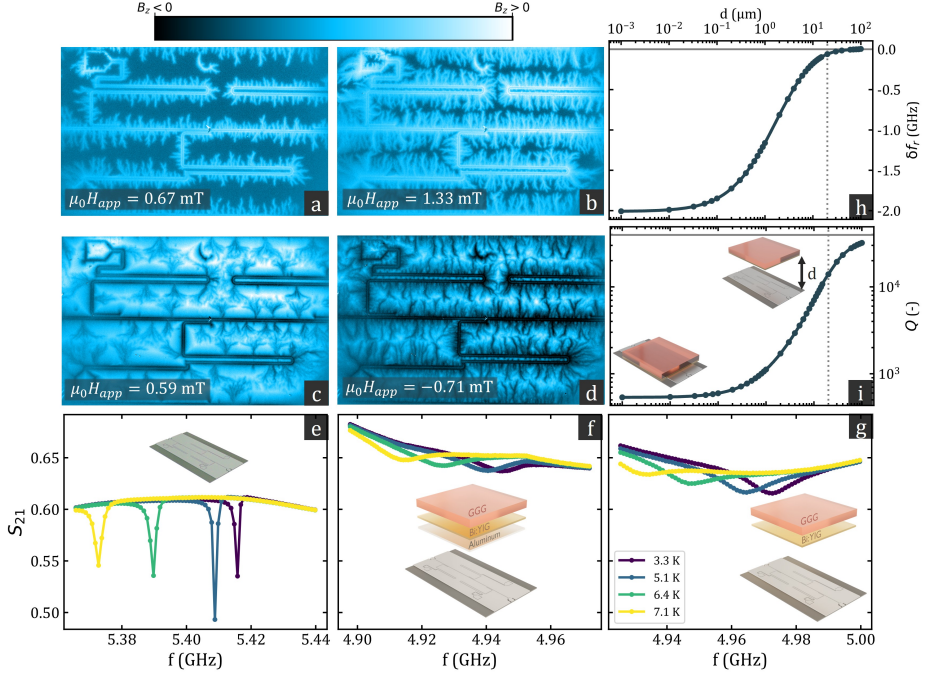


Figure 5.5: (a)–(d) Magnetic flux penetration after zero-field cooling during a field sweep $0 \rightarrow 3.3 \rightarrow -3.3 \rightarrow 0$ mT at $T = 3.2$ K, imaged at 0.67 mT ($\nearrow$), 1.33 mT ($\nearrow$), 0.59 mT ($\searrow$), and -0.71 mT ($\searrow$). (e)–(g) Transmission coefficient S_{21} at -30 dBm and 3.3, 5.1, 6.4, and 7.1 K for the resonator alone (e), with the indicator and its mirror (f), and with the mirrorless indicator (g). (h, i) Evolution of the relative change in resonance frequency, $\delta f_r = f_r(d) - f_r(d = 100 \mu\text{m})$, and quality factor with decreasing distance d between sample and indicator, obtained from Sonnet simulations. The gray line marks the width of the central microstrip line.

deposited on an electrically insulating transparent gadolinium gallium garnet (GGG) substrate with an aluminum mirror at the bottom. This conducting layer absorbs dissipative eddy currents induced by the RF field, severely affecting the resonator performance: the resonance peaks become significantly broader and the resonance frequency decreases (Fig. 5.5(f)). A trade-off is found by removing the metallic mirror layer (Fig. 5.5(g)): the shift in resonance frequency persists, but the resonant lines are better discerned. In addition, because the sample itself is reflective, little loss of MOI signal is observed. Importantly, with the mirror removed, delayed magnetic flux avalanches can be observed even up to a minute after increasing the magnetic field. This phenomenon was not observed with the mirror, because the thin metal layer suppresses or deflects avalanche propagation through eddy-current-induced magnetic braking [148, 149] and its enhanced thermal conductivity facilitates heat dissipation, thereby reducing the thermal feedback loop responsible for triggering additional avalanches. The mirrorless configuration is adopted as the final configuration for simultaneous imaging and transmission measurements. In some cases, the invasive character of the MOI technique can even lead to a complete suppression of the resonance line.

Numerical simulations using the commercial software Sonnet [150], taking into account the exact dimensions of the bilayer indicator ($\varepsilon = 11.99$ for the GGG substrate [151] and $\varepsilon = 10$ for the Bi:YIG [152]), show a significant reduction in both resonance frequency and quality factor as the indicator approaches the sample (Fig. 5.5(h,i)). Notably, the quality factor exhibits reduction even at relatively large distances where the frequency shift remains weak. The resonance frequency only becomes significantly affected at distances of the order of the central track line width ($\sim 20\ \mu\text{m}$). Both parameters reach a plateau at small distances, indicating saturation of the indicator influence in close proximity. The experimentally observed shift of less than 0.5 GHz is consistent with a separation of approximately $5\ \mu\text{m}$, in agreement with typical values reported for MOI setups [75]. The observed downshift can be attributed to the increase in

the effective permittivity of the medium surrounding the coplanar waveguide, since $f_r \propto 1/\sqrt{\epsilon_r}$ [50]. This influence is purely dielectric: the excitation frequency (5–8 GHz) is orders of magnitude above the uniform ferromagnetic-resonance band of Bi:YIG under in-plane magnetic fields of a few mT. These results highlight the critical trade-off: achieving high MOI resolution requires minimizing the sample-to-indicator distance, but this increases the invasiveness of the technique.

5.6.3 Effect of RF excitation on the avalanche activity

We can now address the central question left open by the sequential measurements: does the RF excitation stimulate the nucleation and propagation of magnetic flux avalanches? To this end, the magnetic field is swept in a stepwise fashion over a complete loop ($0 \rightarrow 3.3 \rightarrow -3.3 \rightarrow 0$ mT), and at each field step a frequency sweep around the resonance is performed. Before and after each RF sweep, MOI images are acquired to visualize the avalanches appearing during the RF excitation. To determine whether observed avalanches are RF-induced or simply of delayed nature, a control measurement is performed in which the frequency sweep is replaced by an equivalent waiting time with the RF power turned off. This protocol is summarized in Fig. 5.6(a), and individual avalanches are color-coded with respect to the field value at which they were detected. All avalanches occurring outside the RF sweep/control period are removed by the image subtraction.

A comparison between RF and no-RF measurements shows that more avalanches appear during the RF sweep (Fig. 5.6(b)–(e)). However, increasing the delay time Δt between field steps reduces the number of avalanches, suggesting that a fraction of them is delayed rather than RF-induced. For larger delay times, the avalanche activity becomes too weak for meaningful statistical analysis.

To investigate the effect of increasing RF power, the distribution

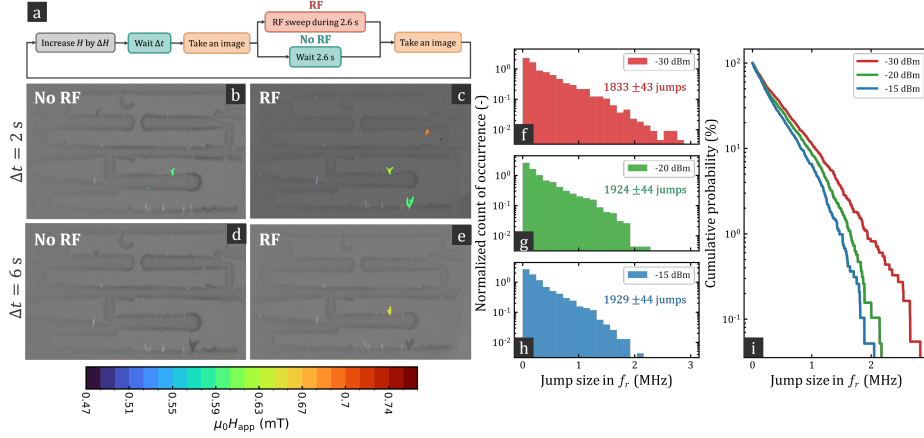


Figure 5.6: (a) Experimental protocol followed to characterize the influence of a RF sweep on magnetic flux avalanches, applied after a zero-field cooling. Subsequent images taken before and after the RF sweep or waiting are subtracted to reveal the avalanches that appeared in between their acquisition. (b)–(e) Avalanche activity appearing after waiting Δt when increasing the magnetic field: (b) control and (c) RF sweep for $\Delta t = 2$ s; (d) control and (e) RF sweep for $\Delta t = 6$ s. (f)–(h) Histograms of the frequency jump sizes $|\delta f_r|$ appearing during a field sweep for -30 , -20 , and -15 dBm; each graph results from 10 field sweeps with a constant field step of 0.0216 mT. (i) Cumulative probability of the frequency jump sizes for the three RF powers.

of resonance frequency jump sizes $|\delta f_r|$ is analyzed for input powers of -30 , -20 , and -15 dBm (Fig. 5.6(f)–(h)), all measured without an indicator. At lower RF power, larger jumps occur, although less frequently: the tail of the histogram extends further into large jump sizes at lower powers, while at higher RF power slightly more jumps are detected but they are smaller. The cumulative probability of jump sizes (Fig. 5.6(i)) shows that the curves overlap for small steps but diverge for large frequency shifts: higher RF powers result in steeper curves, indicating that large jumps are less likely. This behavior can be understood as follows: avalanches result from the decay of a metastable state. The gentler the shaking, the further from equilibrium the metastable state, and therefore the larger the resulting avalanche.

In summary, vortex shaking by the RF field within the low-power, linear Campbell regime ($\mu_0 h_{RF} < H_p$, where H_p denotes the vortex penetration field) slightly enhances magnetic flux avalanche activity, but this remains a marginal effect compared to the impact of changing the applied DC field. Notably, no evidence of excess avalanche activity is found at the positions where the circulating RF current (proportional to the RF field) maximizes. This is consistent with previous reports suggesting that $\mu_0 h_{RF}$ does not need to exceed the vortex penetration field in order to trigger avalanches.

5.6.4 Effect of avalanche activity on the RF transmission

The simultaneous RF-MOI measurements enable the unambiguous association of individual avalanche events with specific jumps in the resonance frequency, extending the analysis beyond what was possible with sequential measurements.

Fig. 5.7(a-c) shows the resonance frequency for the three resonators on the chip during a magnetic field sweep along one of the branches of the $f_r(H)$ loop, as indicated with a dotted box in Fig. 5.4(f). Representative snapshots showing the magnetic activity during the

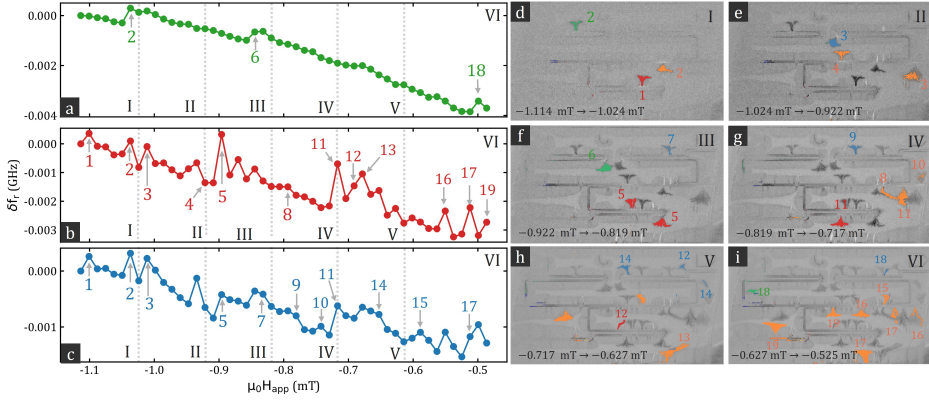


Figure 5.7: (a)–(c) Resonance frequency as a function of applied magnetic field for resonators 1, 2, and 3, respectively, along the branch indicated by the dotted box in Fig. 5.4(f). (d)–(i) MOI images corresponding to the field intervals defined in (a)–(c). Each new avalanche is colored according to its nucleation site: resonator 1 (green), resonator 2 (red), resonator 3 (blue), and feedline or sample edges (orange). Each avalanche event is labeled with a number and linked to the corresponding jump in resonance frequency.

magnetic field intervals, and marked with Roman numerals, are shown in the left panels. In the snapshots, each emerging avalanche is assigned a color code that permits identification of the nucleation sites. Avalanches nucleating in resonator 1 are colored green, on resonator 2 with red, and on resonator 3 with blue. Avalanches forming anywhere else, specifically at the feedline or the sample's edges, are colored in orange. Each avalanche event has been labeled with a number and linked to the corresponding jump in resonance frequency in panels (a-c).

It can be seen that jumps in resonance frequency correlate with magnetic flux avalanches nucleating on the corresponding resonator. Upward jumps can be explained by an effective decrease of the magnetic field at the edges of the patterned sample after an avalanche event. This is so because of the particular magnetic history in which the sample is initially fully penetrated by a positive magnetic field, and avalanches of anti-vortices tend to decrease the local average magnetic field. Since the RF current distribution is inhomogeneous and peaks at the edges [55], a decrease in the local magnetic field also results in a reduction of dissipation caused by vortex shaking. Additionally, a decrease in the magnetic field reduces the kinetic inductance, thereby increasing the resonance frequency. After this event and until the next avalanche occurs, the resonance frequency is observed to decrease as magnetic pressure builds up again. Downward jumps are also observed, and they are actually more frequent and cover a larger magnetic field range than upward jumps. Their origin can be understood considering the fact that an avalanche effectively creates a normal region with a local increase in magnetic flux, and therefore with higher kinetic inductance.

Interestingly, we found that events in one resonator may affect the neighboring resonators. This non-local effect is manifested, for instance, in event 1, which is triggered in resonator 2, and affects the resonance frequency not only of resonator 2, but also that of resonator 3. Similarly, event 2, which involves two simultaneous avalanches, one located in resonator 1 and another along the excitation line, this

event is detected in all three resonators. This cause-effect non-locality is intimately related to the fact that, for thin-film geometry in a perpendicular magnetic field, the magnetic field at a given position is determined by the distribution of the current density over the whole film [37]. The immediate consequence of this observation is that strategies for improving flux pinning need to consider the non-locality of the problem. In other words, reinforcing pinning there where the density of vortices maximizes may only be a partial solution to a mitigation plan.

It may also come as a surprise to see that magnetic flux avalanches are favored along flat borders, rather than corners, where the concerted action of current crowding and nanofabrication defects is expected to facilitate the penetration of vortices. This counterintuitive observation occurs at intermediate temperatures and has been explained in Ref. [124, 153] based on the field dependence of the critical current. Nevertheless, at low temperatures, we have observed that avalanches first appear at the U-shaped turn of the resonators.

One cannot help but note that in resonator 1, avalanches are markedly less frequent than in the other two resonators. We attribute this observation to the fact that this resonator is effectively at a higher temperature. The technical reason is that the size of the chip is slightly larger than the cold finger of the cryostat, making thermalization of this part of the substrate slightly less efficient. Since avalanche activity decreases with increasing temperature, this resonator is expected to exhibit fewer avalanche events.

5.6.5 Electromagnetic simulations

The variability in frequency jump size can arise from multiple factors: the multiplicity of co-occurring events, the size of the avalanche, and its physical location. To disentangle these contributions, method-of-moment simulations of the electromagnetic response are performed using Sonnet, in which a magnetic flux avalanche can be emulated either by a local increase in resistance (R_{RF}) or a local increase in

kinetic inductance (L_k). As shown in Fig. 5.8(a), the former approach decreases f_r and simultaneously reduces Q ; the latter decreases f_r with little change in Q . Since experimentally no net reduction of Q is observed during a single jump event (Fig. 5.4(g)), emulating avalanches as local changes in L_k is the physically appropriate choice. This is physically understood as follows: an avalanche renders a region of the superconductor transiently resistive, but by the time the $S_{21}(f)$ measurement is carried out the region has cooled back below T_c and can be modeled as a region with a *frozen* high local magnetic field, which raises L_k and reduces f_r . Conversely, for anti-vortex avalanches the local field decreases, L_k decreases, and f_r increases.

Using the exact geometry of Device #2 (Fig. 5.8(b)), the simulations compute the resonance frequency shift δf_r as a function of the position of a rectangular avalanche patch (kinetic inductance increased by a factor of ten) along the resonator. The current density distribution at resonance is superposed with the selected 20 possible avalanche locations along both the outer and inner perimeter of the U-turn. We stress that these calculations are static, linear electromagnetic simulations, not intended to model the dynamics of avalanche nucleation or propagation; instead, they provide a static picture of how a completed avalanche modifies the microwave response. Three main conclusions follow.

First, flux avalanches have a higher impact on f_r where the amplitude of the sheet current density is higher. This correlation is made more apparent in panel (d), where δf_r of the 20 possible avalanche locations is plotted as a function of the amplitude of the sheet current density squared. Previous studies have similarly shown that shifts in the resonant frequency can be traced back to the surface magnetic field and, by extension, to the surface current density, thereby confirming that the measured frequency shifts provide insight into the amplitude of the current distribution [92, 98].

Second, avalanches appearing at the inner part of the U-shaped meander seem to have a lower impact than avalanches at the outer perimeter. This could be attributed to the fact that inner avalanches

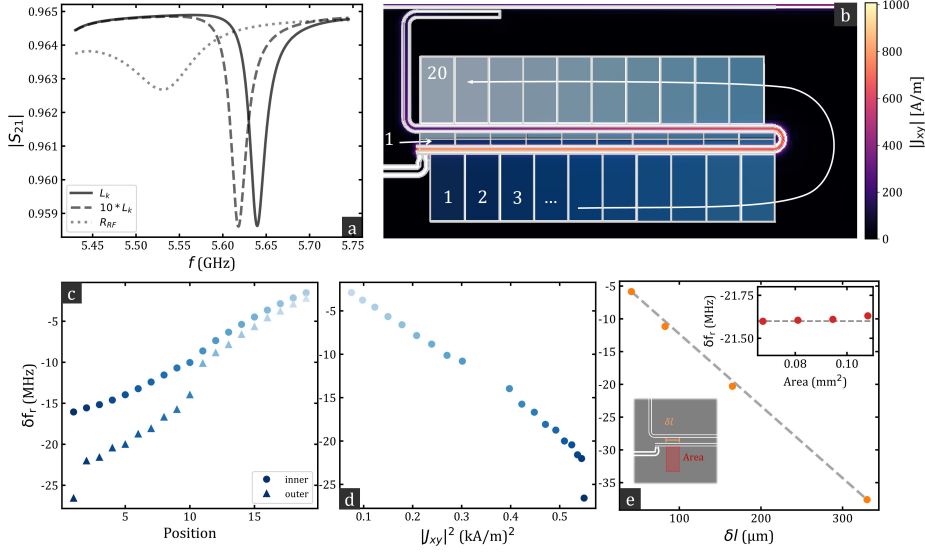


Figure 5.8: (a) Simulated S_{21} spectrum of the resonance peak for a superconducting patch modeled with the same kinetic inductance as the rest of the material, an increased kinetic inductance ($10 L_k$), and a resistive component (R_{RF}). (b) Simulated current density at resonance for an applied power of +13 dBm, superposed with the 20 outer and 20 inner avalanche locations considered. (c) Simulated relative change in resonance frequency as a function of the avalanche position; triangles correspond to the outer perimeter and dots to the inner perimeter, with marker colors encoding the location. (d) Change in resonance frequency as a function of the sheet current density squared at the center conductor near the simulated patch. (e) Simulated relative change in resonance frequency as a function of the contact length with the conductor edge; the upper inset shows δf_r versus the patch area, while the lower inset illustrates the contact length and area of the patch.

cover a substantially smaller surface than outer ones. However, we have verified that only the contact length matters (see panel (e)), not the surface of the avalanche (see inset of panel (e)). Thus, one is tempted to conclude that avalanches at the inner border of the resonator are less harmful than those at the outer border. However, care must be exercised with this interpretation, since the simulations incorporate the effect of the magnetic field through the local changes in L_k , which need to be adapted accordingly for the inner and outer rims of the resonator. Indeed, the field-focusing effect at the inner border should lead to a proportionally larger L_k for inner avalanches compared to outer ones.

Third, we observe that avalanches occurring at one resonator do not influence the response of the neighboring resonators. This absence of inter-resonator coupling in the simulations is an expected consequence of the model's design and does not contradict the experimental observations. Experimentally, non-local inter-resonator effects arise because, in a thin-film superconductor under a perpendicular field, the current density at any point is determined by an integral over the current distribution of the entire film [154]: a sudden vortex rearrangement in one region necessarily modifies screening currents throughout the chip, including in neighboring resonators that did not host the avalanche. This is a fundamentally non-linear and non-local process [155]. The Sonnet model, which represents avalanches as static, local perturbations of kinetic inductance on a linear perfect conductor, cannot capture this dynamical current redistribution. The model is accurate for its intended purpose (quantifying the static electromagnetic response to a completed, localized avalanche), but its linearity and locality preclude the reproduction of inter-resonator effects. Capturing this non-local coupling self-consistently would require thermomagnetic models that co-evolve the vortex distribution, local temperature, and current density in time [34, 155]. Furthermore, we do not observe an influence of avalanches at the feedline on the resonator.

5.6.6 Dynamic regime analysis

A concluding discussion of the dynamic regimes explored in this study is warranted. It is important to acknowledge that the presence of vortices in superconducting resonators is essentially unavoidable. Most resonators are designed to operate in the gigahertz range, requiring fabrication on a scale of several millimeters. For NbTiN, the superconducting characteristic lengths are $\xi \sim 4$ nm and $\lambda \sim 400$ nm, meaning there is ample space to host a large number of vortices [12]. Since RF currents are localized along the edges, vortices in the perimeter of the structure are particularly deleterious.

Concerning vortex trapping, it has been shown that in a 20- μ m-wide NbTiN strip under field-cooling conditions, trapping can occur at magnetic fields as low as 10 μ T [12]. Taking $B_{c1} \approx 4$ mT and the demagnetization factor $\sim \sqrt{W/t}$ with $W = 5$ mm and $t = 100$ nm, the first penetration field into the outer rim of the sample is estimated as $B_p \sim 20$ μ T, assuming Bean-Livingston barriers are suppressed by border defects and inhomogeneities [153]. The first discontinuity in the resonance frequency is observed at much higher fields (100–500 μ T), corresponding to the magnetic perforation of the outer rim and sudden injection of magnetic flux into the gaps of the structure. This magnetic field is then locally increased by a factor of ~ 10 due to flux focusing [23, 147], rapidly populating the borders of feedline, resonator, and ground plane with vortices.

A much larger length scale ℓ^* dominates the physics of avalanche formation. This length represents the threshold penetration depth above which thermomagnetic instabilities may develop [133, 156]. In Device #2, avalanches were clearly observed within the 60- μ m-wide spacing between the two branches of the U-shaped resonators, as well as in both the central superconducting line of the resonators and the feedline.

Concerning the RF excitation, the presented measurements span a power range of -30 to 0 dBm, corresponding to circulating current densities $J_{RF} \sim 0.1\text{--}3 \times 10^5$ A/cm² and RF fields $\mu_0 h_{RF} \sim 0.03\text{--}$

1 mT after accounting for the flux focusing effect [147]. These values are respectively below the critical current density of NbTiN ($J_c \sim 10^7$ A/cm² [12]) and the lower critical field, confirming that the RF excitation is limited to vortex shaking rather than vortex reconfiguration, i.e., within the linear Campbell regime [146]. This is consistent with the threshold $\mu_0 h_{RF} \sim 5$ mT for non-linear response in NbN films [144]. It is worth pointing out that the magnetic flux avalanches develop primarily in the ground plane or in the spacing between the two branches of the U-shaped resonators, where even lower current densities are expected [55]. Indeed, the edge currents in the ground plane next to the excitation line or the resonators are approximately a factor $(1 + 2g/w) \sim 2$ smaller than in the central line [98].

5.7 Conclusions

We have explored NbTiN CPW resonator structures using magneto-optical imaging and RF transmission measurements, both in sequential and simultaneous configurations. A clear and unambiguous correspondence between the two independent methods is demonstrated.

The sequential measurements allow a direct link between the complex magnetic field penetration and the changes in resonance frequency. Four distinct regimes are identified: a magnetic field screening regime, a magnetic perforation event, a flux avalanche regime, and a smooth flux removal regime. Thermomagnetic instabilities, active below $T \lesssim T_c/2$, drastically impact the resonance frequency through sudden jumps in both f_r and Q_i . The threshold field for avalanche activity is consistently extracted from both MOI and RF measurements, demonstrating that superconducting resonators are a valuable non-invasive tool to investigate magnetic flux dynamics. A roadmap for improving field resilience is proposed: reducing the ground plane width, reducing the resonator line width, and carefully considering the effect of artificial pinning sites at low temperatures.

The simultaneous measurements, enabled by the setup upgrade, extend these findings in two key ways. First, they establish that RF excitation in the linear Campbell regime has only a marginal effect on avalanche activity, providing a direct experimental answer to a long-standing open question. The increased avalanche activity under RF excitation is real but small, and correlates with the degree of metastability rather than with the local amplitude of the RF current. Second, the simultaneous measurements enable unambiguous tagging of individual avalanche events to specific jumps in resonance frequency, revealing the non-local character of the electrodynamics: avalanches in one resonator, or even in the feedline, can affect all resonators on the chip. Electromagnetic simulations confirm that the most deleterious avalanche locations are those with the highest RF current density and that the contact length of the avalanche with the conductor edge more than its area determines its impact on f_r .

The invasive character of the MOI technique is quantified and shown to shift the resonance frequency and reduce the quality factor due to increased effective permittivity and dielectric losses. Less intrusive techniques, such as Quantum Diamond Microscopy [157], could be implemented in future work to make the reduction of the threshold field for triggering magnetic flux avalanches under RF excitation more apparent. Finally, it is worth pointing out that the magnetic flux avalanche regime is ubiquitous in thin films made of conventional superconductors and may affect most technologically mature quantum devices. High- T_c superconducting thin films, which are less susceptible to these events, may offer an alternative route to magnetic-field-resilient resonators.

Chapter 6

Scanning Laser Microscopy of High- T_c $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ Devices

The previous chapter established that the erratic field response of a superconducting resonator is not simply experimental noise but the consequence of thermomagnetic flux avalanches, and that the magnitude of the frequency jump that each avalanche produces depends on *where* the avalanche is triggered. That conclusion leaves two threads hanging. The first concerns the material: flux avalanches are notoriously hard to trigger in YBCO [133, 156], which therefore offers a resonator platform free of the erratic field response reported in Chapter 5. The second thread is the more pressing one, and it concerns the second question posed in the Introduction: if the harmfulness of an avalanche is set by its location, then one must know which regions of a resonator are actually more susceptible to affect its resonance frequency. Chapter 5 mapped the magnetic flux density and the associated screening currents. However, no information is provided about the radio-frequency current of the resonant mode, which is a distribution of an entirely different nature, set by the circuit geometry and the standing-wave pattern rather than by the applied

field, and which is present even in zero magnetic field. Mapping it requires a different microscope.

This chapter takes up both threads, on $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ (YBCO) coplanar half-wavelength resonators, and with scanning laser microscopy (SLM), the second of the two local probes on which this thesis rests. The strategy remains the one announced in the Introduction: map, point by point, the local quantity that governs the phenomenon at hand, and confront that map with the averaged microwave observable that a conventional measurement would return. What changes is the local quantity, now the RF current density, and the means of accessing it, a focused laser beam that perturbs one micrometer-sized spot at a time while the resonance is being tracked.

A note on the structure of this chapter. The results presented here were obtained over two experimental investigations, separated in time. Sec. 6.1, adapted from N. Lejeune *et al.*, “Spatially resolved photodoping of oxygen-deficient $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ resonators” [47], relies on a simple home-built laser-scanning arrangement combining a widefield illumination path with a galvanometer-steered laser, which is introduced in the section in which it is used. Sec. 6.2, adapted from N. Lejeune, S. Van der Heyde and A. V. Silhanek, “Time-Resolved Thermal Susceptibility Mapping via Low-Temperature Scanning Laser Microscopy” [158], instead employs the new dedicated scanning laser microscope whose optical layout and the calibration of its optical and thermal spot size have already been already presented in Chapter 4.

Personal contribution. In the photodoping study of Sec. 6.1, the YBCO films were grown and patterned at ICMAB-CSIC; I led all the remaining aspects of that work, namely the measurements, the phenomenological convolution model, and the analysis. The scanning laser microscope used in Sec. 6.2 and described in Chapter 4 was specified, designed and assembled by myself, from the selection of the individual components to the control and acquisition software. Its optical and thermal spot sizes were calibrated together with S. Van der

Heyde, while the time-resolved measurements, the COMSOL thermal simulations underpinning their interpretation, and the analysis of the delayed response were led by myself.

6.1 Photoexcitation of YBCO resonators

6.1.1 Introduction

When considering the interaction of electromagnetic waves with conventional superconductors, one encounters a large diversity of physical mechanisms and phenomena, including high reflectivity [159], inverse Faraday effect [160], light-induced superconductivity [161–163], quasiparticle poisoning [95], bolometric (thermal) effect [90, 92], room-temperature transient superconductivity [164], photovoltaic effect [165] and many more. Cuprate superconductors behave the same way under short illumination pulses and for low-photon doses. However, some cuprates, such as oxygen-depleted $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ (YBCO), exhibit a striking enhancement of their superconducting properties at much longer illumination times [166]. More concretely, the superconducting transition temperature increases considerably, and the photo-induced change remains after the light has been switched off as long as the sample is kept below a certain temperature ($T^* \sim 250\text{ K}$). Besides this property, coined persistent photoinduced superconductivity (PPS), there is also an increase of conductivity in the normal state, known as persistent photoinduced conductivity (PPC). For temperatures exceeding T^* , the system relaxes very slowly to its initial conditions. As for the underlying mechanism, after much debate, there seems to be some consensus on the fact PPS and PPC are associated with the presence of oxygen vacancies in the basal CuO chains. Recently, it has been proposed that PPS results from photo-induced oxygen ordering within the CuO chains, leading to a decrease in the electronic scattering rate, whereas PPC can be naturally explained by the excitation of electron-hole pairs followed by electron trapping in the oxygen vacancies and the promotion of holes

to the conduction band (photodoping) [167]. This level of control has enabled practical implementations, such as writing Josephson junctions using a scanning near-field microscope [168], demonstrating the method's potential for device engineering.

A particularly suitable way to investigate the photoresponse of YBCO superconducting systems is through coplanar transmission line resonators. These circuits can achieve quality factors as high as $Q \approx 10^7$ in low critical temperature superconductors [1, 169] and on the order of $10^3 - 10^4$ in cuprate superconductors [170–172], offering unparalleled sensitivity to minute changes induced by external stimuli or environmental variations [138]. The demonstration of bolometric changes under laser light on superconducting thin films has spurred the use of localized photoresponse as a probe of superconducting resonators by scanning a laser spot across the transmission line. In this technique, an amplitude-modulated laser beam is focused onto the surface of the superconductor, and the absorption of this light results in a modulation of the local temperature. Since the superconducting penetration depth is dependent on temperature, the local penetration depth and thus the local stored magnetic field energy will be modulated [39]. In other words, the effect of illumination is attributed to the photoinduced change of the kinetic inductance due to laser heating. This approach has permitted imaging the local microwave current density in a superconductor [39, 40] and enables tuning the resonant frequency of the resonator with negligible change of the quality factor [173].

The investigations described above were conducted on optimally doped YBCO thin-film resonators, where the low density of oxygen vacancies renders persistent photo-induced effects negligible, making thermal heating the dominant mechanism. In this Chapter, we explore the long-term photo-response of oxygen-depleted YBCO half-wavelength resonators, focusing on the effects of photodoping. This mechanism initially counteracts the bolometric effect and eventually becomes dominant, exhibiting persistent and metastable behavior. By locally illuminating a spot of comparable size as the width of the

resonator's track, at the center of the resonator (where the radio-frequency field is strongest), we are able to enhance the resonator's characteristics. This improvement has been previously demonstrated through the introduction of non-uniform columnar defects created by heavy-ion irradiation [41]. Our approach is reversible and enables repeated and multiple comparisons using the same device. This method is applied to map the standing waves at both the fundamental resonance and the second harmonic and complements the bolometric method described in Ref. [92, 98, 99, 174].

6.1.2 Sample fabrication and resonator design

Epitaxial $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ films of thickness 100 nm were grown on top of $\text{La}_{0.3}\text{Sr}_{0.7}(\text{Al}_{0.65}\text{Ta}_{0.35})\text{O}_3$ (LSAT) single crystal substrates by pulsed laser deposition. LSAT is a convenient choice due to its low dielectric constant, structural stability, low losses, isotropic microwave response, and it has no twinning [99, 170]. Substrates were heated up to 810 °C, with an oxygen partial pressure of 0.3 mbar and a target-to-substrate distance of 52.5 mm. Ablation of the target material was achieved by using a high fluence laser operating at 2 J/cm² and a frequency of 5 Hz. During the cooling process to room temperature, the samples were maintained in an oxygen atmosphere of 1 bar. This way, an optimally doped film $\delta \approx 0.1$ with superconducting critical temperature $T_c \approx 90$ K was obtained. A second film was deoxygenated through an *ex situ* post-annealing process at 400 °C for 4 hours in a 99.5% N₂ atmosphere, leading to a lower oxygen concentration $\delta \approx 0.7$ with superconducting critical temperature $T_c \approx 60$ K [175]. Prior to sample patterning, the T_c of the films was confirmed using a SQUID magnetometer. Both films were patterned by photolithography and wet etching into 8 mm-long 50Ω $\lambda/2$ resonators with a 20 μm -wide coupling finger-like capacitor, as illustrated in Fig. 6.1(a). The coupling capacitors have been deliberately designed within the weak coupling regime to ensure a quality factor high enough to detect minor changes in the resonance frequency (see Box 2.4) [171]. The

resonance frequency of such a device is set by its geometric and kinetic inductance per unit length through Eq. (2.23), the kinetic contribution following the Gorter–Casimir temperature dependence of Eq. (2.25), as detailed in Sec. 2.1.

The 8 mm-long resonator was designed such that its resonance frequency is close to 5 GHz, with design values $C' = 214.1$ pF/m and $L'_g = 634$ nH/m. The transmission scattering matrix element $S_{21}(f)$ in the vicinity of the resonant frequency is measured through a vector-network analyzer (Rohde & Schwarz ZNL20). A representative example of $S_{21}(f)$ is shown in Fig. 6.1(c). Its asymmetric shape reveals an impedance mismatch [61] and is fitted with Eq. (2.15) of Sec. 2.1.3, which returns the resonance frequency f_0 together with the internal and external quality factors Q and Q_e . For the particular example of Fig. 6.1(c), the fitted resonance frequency is $f_0 = 5.153$ GHz ± 2.5 kHz and the quality factor values are $Q = 2128 \pm 4$ and $Q_e = 8.284 \cdot 10^6 \pm 16 \cdot 10^3$, confirming that the resonator is strongly under coupled and that losses are dominated by internal losses. Fig. 6.1(b) shows the evolution of f_0 with temperature for both investigated YBCO resonators. The observed sharp decrease of f_0 as temperature increases is due to the divergence of the kinetic inductance when approaching the superconducting transition. It is precisely this strong temperature dependence that has motivated the use of these devices as sensitive bolometers. In addition, the resonance frequency is also dependent on the magnetic field and magnetic history, as shown in Sec. 6.1.5 [172]. Fig. 6.1(d) shows the magnetic field landscape obtained by the magneto-optical Faraday rotation technique [75] (see Chapter 3) of the optimally oxygenated resonator at $T = 3.2$ K and under an out-of-plane magnetic field $\mu_0 H_a = 0.6$ mT after zero-field cooling. This technique enables probing the degree of uniformity in the superconducting properties, as well as identifying the entry points of magnetic flux.

The experimental setup used to investigate the photoresponse of the planar resonator is shown schematically in Fig. 6.1(a). The sample is mounted inside a closed-cycle optical cryostat, allowing

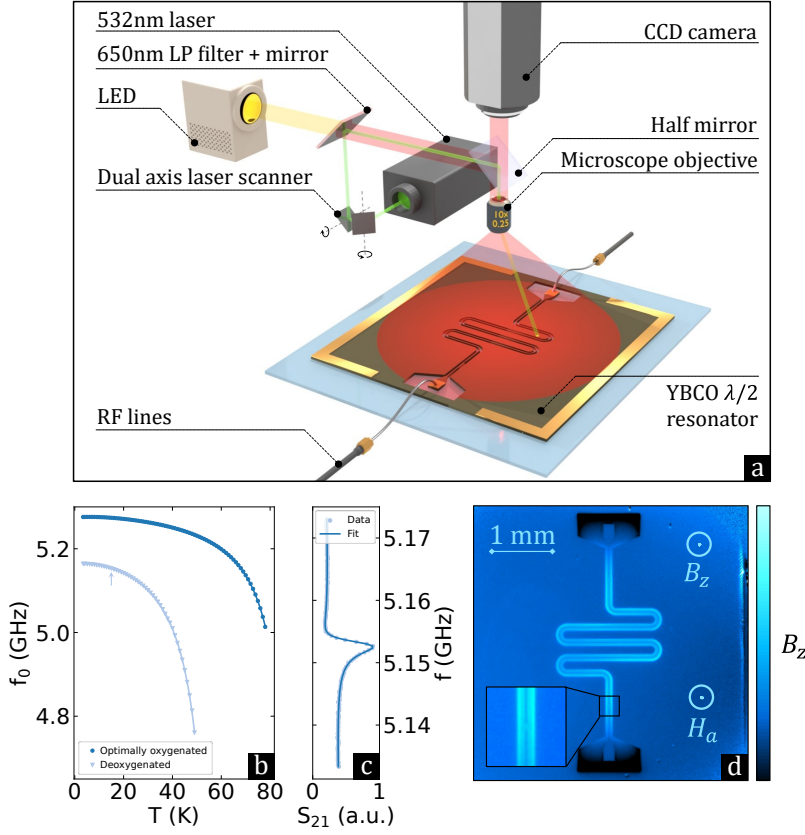


Figure 6.1: (a) Schematic representation of the experimental setup that features a double illumination composed of a widefield red-filtered LED used to image the resonator concurrently with a 532 nm laser that is focused on a desired zone of the sample. When photoexciting the entire sample, the 650 nm long-pass (LP) filter and mirror are replaced by a 550 nm bandpass filter, and the laser is not used. (b) Resonance frequency for the optimally oxygenated YBCO resonator and the deoxygenated one. The resonance frequency drops at temperatures close to T_c due to the diverging kinetic inductance. (c) Selected spectrum measured at 15 K on the deoxygenated sample. By fitting the transmission coefficient $S_{21}(f)$ with Eq. (2.15), it is possible to obtain the resonance frequency for all explored temperatures. At temperatures close to T_c , the uncertainty on the resonance frequency rapidly increases because of the decrease in the intrinsic quality factor. (d) Magneto-optical image of the zero-field-cooled optimally oxygenated resonator at $T = 3.2$ K and applied an out-of-plane field $\mu_0 H_a = 0.6$ mT. Black color corresponds to local magnetic field with out-of-plane component opposite to the applied field. The inset shows a zoom of the coupling capacitor region, evidenced by the bright spot at the center.

for simultaneous RF measurements, optical imaging, and optical excitation. Two illumination modes are employed in parallel. In the widefield illumination mode, a white LED provides broadband light that passes through an infrared blocking filter and a 550 nm bandpass filter with a 100 nm bandwidth. This filtered light uniformly illuminates the sample over its entire area. The resulting images are captured using a CCD camera.

In the localized photoexcitation mode, a 532 nm laser beam is steered by a dual mirror scanner and then passes through a carousel of neutral density filters used to modulate its power. The beam is subsequently reflected toward the sample by a 650 nm long-pass filter that features a mirror coating on the side facing the laser. This element reflects the laser beam while also transmitting about 1% of the LED light (used for widefield imaging) and filtering its spectrum, thereby combining both light paths without interfering with the laser. The reflected laser beam is then focused onto the sample through a $10\times$ microscope objective, producing a spot approximately 60 μm in diameter. Even though the photoexcitation efficiency is stronger in the red part of the visible spectrum [176], the photon dose per unit time coming from the laser is about 10^6 times larger than from the red broadband illumination, ensuring that all photoexcitation effects detected are caused by the laser. Images are recorded with the laser beam blanked, using the CCD camera with exposure time adjusted to capture the faint red illumination (see Sec. 6.1.5).

6.1.3 Results and discussion

Uniform photon irradiation

The critical temperature for both samples can be extracted from the temperature dependence of f_0 shown in Fig. 6.1(b) by fitting with Eqs. (2.23) and (2.25). This yields, respectively for the optimally oxygenated and deoxygenated samples, a similar geometrical inductance for both devices of $636.8 \pm 0.5 \text{ nH/m}$ and $633.9 \pm 1.3 \text{ nH/m}$, in good agreement with the theoretical design value (634 nH/m cal-

culated using elliptical integrals for CPW structures (see Box 2.2), and a critical temperature $T_c = 90.0 \pm 0.1$ K and $T_c = 58.1 \pm 0.1$ K. The zero-temperature kinetic inductance of the optimally oxygenated sample $L'_{k,0} = 18.1 \pm 0.5$ nH/m is significantly lower than for the deoxygenated sample $L'_{k,0} = 49.9 \pm 1.4$ nH/m, manifesting that the latter has a larger sheet resistance and lower superconducting gap.

After characterizing the properties of the deoxygenated sample, the transmission $S_{21}(f)$ was recorded over time every 10 seconds, and the green-filtered LED lamp was switched on for two minutes after the start of the experiment, irradiating the sample with 3.6 mW of 550 nm light. Fig. 6.2(a) shows the time evolution of f_0 at fixed temperature $T = 30$ K $\approx T_c/2$. The dose applied over the entire illumination time is estimated to $D = 5 \cdot 10^{21}$ photons/cm². As soon as the illumination is switched on, a sharp drop of f_0 is observed (see inset of Fig. 6.2(a)), as a result of the heating caused by the light absorption. The thermally-induced drop takes place on a short timescale of the order of hundreds of milliseconds (see Sec. 6.1.5), and is followed by a recovery with a characteristic time scale of the order of hours, resulting from the photoinduced superconductivity. In this regime, $f_0(t)$ follows a stretched exponential functionality in accordance with previous reports [176, 177], such that the global response of the resonance frequency can be modeled using

$$f_0(t) = f_{0,T} - \delta f_{thermal} + \delta f_{PE} \left(1 - \exp \left\{ \left(- \left(\frac{t - t_0}{\tau} \right)^\beta \right) \right\} \right), \quad (6.1)$$

where $f_{0,T}$ is the dark resonance frequency at temperature T , $\delta f_{thermal}$ is the drop in resonance frequency due to light-induced heating, δf_{PE} is the saturation increment of the resonance frequency due to photoexcitation, $t_0 = 2$ min is the time at which the illumination starts, and β is a dispersion parameter. The fit of the experimental data, shown in Fig. 6.2(a), yields $f_{0,30K} = 5.103$ GHz, $\delta f_{thermal} = 3.77 \pm 0.02$ MHz, $\delta f_{PE} = 14.04 \pm 0.1$ MHz, $\tau = 53 \pm 1.1$ hours and $\beta = 0.497 \pm 0.002$.

From the thermally-induced frequency drop, one can infer the temperature load using the previously acquired slope of $f_0(T)$ (Fig. 6.1(b)) yielding a global temperature rise in the superconducting film of about 0.65 K. Furthermore, incorporating the saturation frequency $f_{sat} = f_0 + \delta f_{PE}$, and the obtained geometrical inductance, into Eq. (2.23), the kinetic inductance at saturation can be estimated as 62.1 nH/m compared to the initial value of 65.75 nH/m at 30 K. The fact that reaching 95% of the asymptotic state requires continuously photoexciting the sample for approximately 9τ , which in the case of the LED source is about 20 days, makes it highly time consuming to reach the saturation frequency. Repeating the experiment with a higher-power laser source permits reducing the time required to reach saturation without changing the asymptotic value, as the PPS is only a function of the absorbed dose. It is worth pointing out that, unlike most previous investigations of the PPC and PPS based on $R(T)$ measurements, in which the change in T_c or the normal state resistivity is reported, in this work we directly probe the changes in superconducting properties while at sub-critical temperatures.

Persistent photoinduced superconductivity manifests itself as an increase of the superconducting critical temperature. This effect should become visible as the $f_0(T)$ curve shifts toward higher temperatures. In addition, there is also an increase of conductivity in the normal state (persistent photoinduced conductivity), which in turns reduces the zero-temperature kinetic inductance ($L'_{k,0} = \frac{1}{w} \frac{\rho \hbar}{\pi \Delta(0)}$). A manifestation of both these effects is shown in Fig. 6.2(b), in which the temperature dependence of the resonance frequency $f_0(T)$ is measured after 42-hour-long illumination and compared with the response before illumination. The increase in resonance frequency δf_0 resulting from the photoexcitation is shown in Fig. 6.2(c).

Focused photon irradiation

Once we have characterized the influence of uniform photon irradiation on the deoxygenated YBCO resonator, we proceed to implement

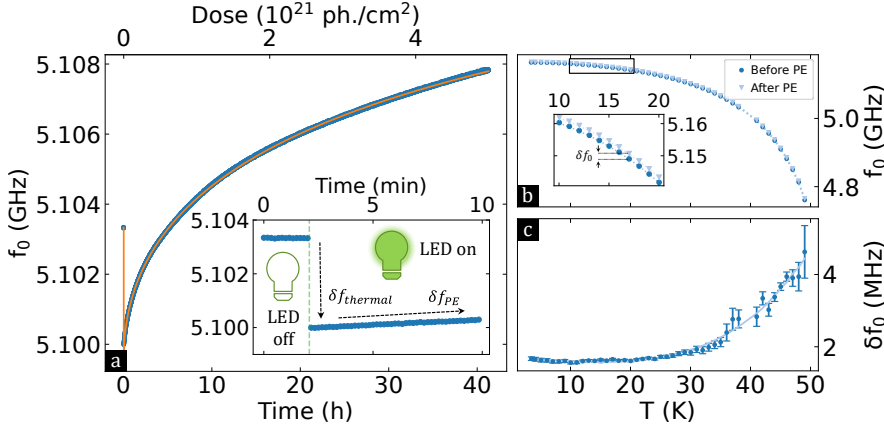


Figure 6.2: (a) Evolution of the resonance frequency of the deoxygenated YBCO resonator with time. When a green LED illuminates the sample, starting at $t = 2$ min, the resonance frequency suddenly drops due to the light absorption-induced heating. On a much slower timescale, the resonance frequency increases due to the photoexcitation. (b) Temperature dependence of the dark resonance frequency of the deoxygenated sample before and after illuminating the entire sample with a dose of $D = 5 \cdot 10^{21}$ photons/cm². The inset zooms in on the upshift in resonance frequency δf_0 . The dashed lines correspond to fits of Eq. (2.23) to the data points. (c) Frequency shifts extracted from panel (b). For temperatures lower than $T_c/2$, the upshift is constant and approximately 2 MHz, then for higher temperatures it diverges. The solid line corresponds to the frequency shift predicted by the fits in panel (b).

focused photon irradiation as a local probe, allowing us to map the sensitivity of the device under local excitations. To that end, we focus a laser beam into a $60\text{ }\mu\text{m}$ -diameter spot, which is raster scanned throughout the resonator's surface while tracking the changes in f_0 . As a consequence of the large difference in timescales of the bolometric effect and the photoexcitation, they can be discriminated using the protocol illustrated in Fig. 6.3(a). For each laser position, the resonance frequency $f_{0,dark}^{(i)}$ is recorded with the laser blanked. Then, the laser beam is unblanked, and another resonance frequency measurement $f_{0,light}$ is acquired. Eventually, after 60 seconds of irradiation, the laser is blanked again, and the resonance frequency $f_{0,dark}^{(f)}$ is probed one last time. The result of the three-point measurements is shown in Fig. 6.3(b), and the inset emphasizes the slight shift toward higher frequencies of the resonance peak after the photoexcitation. The negative frequency shift, purely due to heating, is given by $\delta f_{thermal} = f_{0,light} - f_{0,dark}^{(i)}$, whereas the photoexcitation-induced shift to higher frequency is given by $\delta f_{PE} = f_{0,dark}^{(f)} - f_{0,dark}^{(i)}$. As shown in Fig. 6.2(a), the time required for the PE to compensate the bolometric effect is of the order of hours, and hence, the PE shift is expected to be orders of magnitude smaller than the one obtained from bolometric effects. On top of the purely electrically probed change, the reflectivity of the sample surface, measured locally using the CCD camera, increases over the same timescale, as discussed in Sec. 6.1.5.

The resulting thermal and photoexcitation maps ($\delta f_{thermal}(x, y)$ and $\delta f_{PE}(x, y)$) on the deoxygenated sample, probed at the fundamental frequency for a scanning grid of 40×50 pixels over a $2230 \times 2790\text{ }\mu\text{m}^2$ area centered on the resonator, are shown in Fig. 6.3(c,d). A clearly different response is observed through these two methods. Indeed, the thermal response exhibits a more pronounced frequency change in the ground plane, in particular around the mid-point of the resonator, than along the central track and the gap of the resonator. In addition, the excitation and detection lines, capacitively coupled to the resonator, are visible. In contrast to that, the photoexcitation map

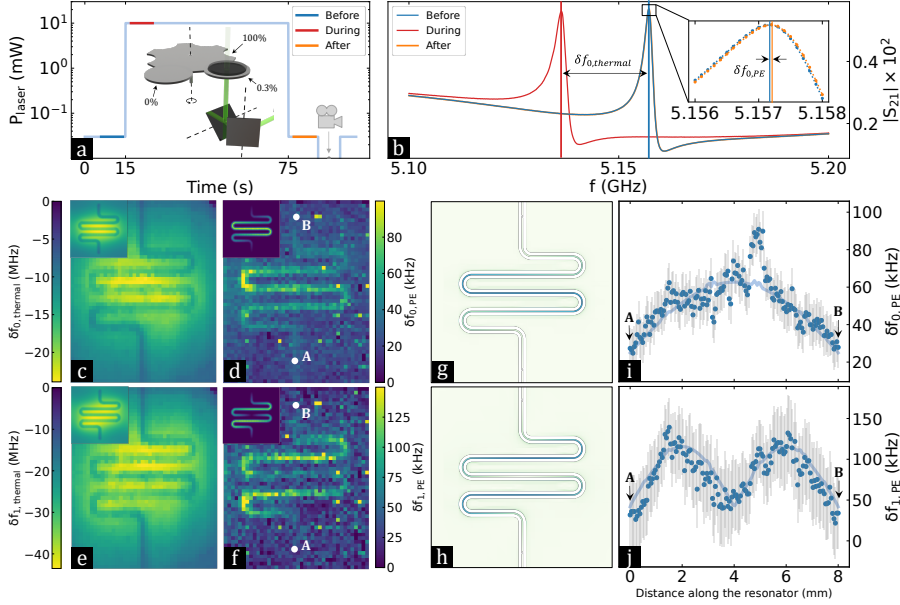


Figure 6.3: (a) Schematic representation of the protocol executed at each laser position. Three measurements are performed and synchronized with the unblanking of the laser. At each position, the sample is locally photoexcited for 60 s. A picture is taken using a CCD camera with the laser fully blanked to keep track of the change in reflectivity. The inset shows a schematic view of the neutral density filters carousel used to modulate the laser power by respectively blanking (0%), strongly attenuating (0.3%) or exposing the sample (100%). (b) $S_{21}(f)$ spectra before, during and after the irradiation. The inset is a zoom emphasizing the slight shift of the resonance peak toward higher frequencies after the irradiation. (c) Thermally-induced ($\delta f_{0,\text{thermal}}$) and (d) photoexcitation-induced ($\delta f_{0,\text{PE}}$) resonance shift from the fundamental resonance acquired while sweeping the laser over the sample. (e) and (f) are the same maps as panels (c) and (d), but for the second harmonic at double the fundamental frequency. The insets in panels (c-f) show the resulting map obtained using the convolution model described in the main text. (g) (resp. (h)) corresponds to the simulated current density at the fundamental (resp. second harmonic) resonance. (i) and (j) are line profiles following the resonator's shape of the PE shifts for the fundamental and the second harmonic, respectively. The path direction starts at point A and ends at point B as indicated in the panels (d) and (f). The gray bars correspond to uncertainties in the measured shifts. The continuous blue lines represent the current density extracted from panels (g) and (h).

reveals essentially the central track of the resonator, with the ground plane and excitation/detection feedlines remaining largely insensitive. A similar trend is obtained at the second harmonic $f_1 = 2f_0$ ($l = \lambda$, where λ is the wavelength associated with the resonance frequency), as shown in Fig. 6.3(e,f).

These findings suggest that the thermal response is much broader than the laser spot size due to the heat diffusion in the sample and substrate, whereas the photoexcitation induced directly by the irradiation remains localized to the optical spot. Indeed, numerical modeling reveals that the local temperature at the center of the laser spot increases by approximately 5 K and the thermal spread, characterized by the full width at half-maximum, reaches around 120 μm in diameter. Notably, the initially sharp Gaussian heat source broadens significantly in the thermal response, leading to an increase in temperature of more than 0.5 K across a region 500 μm in diameter, covering roughly one-third of the resonator's surface. The spatial temperature distribution and more details about the numerical calculations are provided in Sec. 6.1.5.

Phenomenological model

A simple model can account for the observed contrasting behavior between bolometric and photoexcitation maps. Indeed, the convolution of physically grounded kernels, i.e. kernels able to describe the mechanisms at play, can faithfully capture the obtained maps. A graphical representation of the phenomenological model described below is presented in Sec. 6.1.5.

The first step consists of identifying the absorption function M_{abs} . For the resonance frequency to be impacted by the irradiation, the laser beam must be absorbed by the sample. Note that the laser irradiation is more absorbed by the superconducting YBCO sample (86% transparent) [178] than by the LSAT substrate (20% transparent) [179]. Therefore, when the laser beam is parked in the gap of the CPW geometry, the sample is substantially less heated, and the

change in frequency is attenuated. In contrast, when it sits in the ground plane, the frequency change is accentuated. This is modeled by taking the convolution of the sample local absorbance $S(x, y)$ with the Gaussian profile from the laser $L(x, y)$. The end result is a qualitative approximation of the quantity of light absorbed at each scanned pixel by the laser, $M_{abs}(x, y) = S(x, y) * L(x, y)$.

Second, we convolve the theoretical current density of the sample $J(x, y)$ with a kernel representing either the bolometric ($G_R(x, y)$) or the photoexcitation effect ($G_r(x, y)$). Here, the argument justifying the use of the current density lies in the fact that measurements are sensitive to changes in L'_k , which in turn is correlated to the current density ($L'_k \propto I^2$) [53, 64]. Previous studies have similarly shown that shifts in the resonant frequency can be traced back to the surface magnetic field and, by extension, to the surface current density, thereby confirming that the measured frequency shifts provide insight into the underlying current distribution [98]. To that end, the current density along the resonator's track line is approximated by a sine-like profile with one anti-node for the fundamental mode, and two for the second harmonic. We assume that the current density at resonance is negligible in the rest of the ground plane, which is consistent with simulations. In the case of photoexcitation, since the physical mechanism implies direct absorption of photons without diffusion, the convolution uses a Gaussian profile $G_r(x, y)$ of the same size as the laser optical spot. In contrast to that, for the bolometric response, the kernel corresponds to a function $G_R(x, y)$ reproducing the temperature spread resulting from the finite element simulation. Ultimately, the element-wise products, denoted by $\odot$, given by $M_{abs}(x, y) \odot [J(x, y) * G_r(x, y)]$ and $M_{abs}(x, y) \odot [J(x, y) * G_R(x, y)]$ express the requirement that, for any frequency modification to occur, two conditions must be met simultaneously at each point (x, y) : (1) light must be absorbed (as indicated by M_{abs}), and (2) the resonator must be sensitive at that location (as described by the convolution $J * G$). The results of these models are presented in the insets of Fig. 6.3 (c-f) and qualitatively reproduce the measured maps. The predicted bolometric response

even captures the observed broadening of the thermally induced frequency shift region over the whole resonator at $2f_0$, resulting from the fact that the current density is delocalized due to the presence of two anti-nodes in the standing wave. A comparison of the profiles obtained from the experimental data and the phenomenological model can be found in Sec. 6.1.5.

It is worth mentioning that the bolometric response has been extensively investigated as a method to map current distribution in superconducting waveguides and resonators [98, 174]. In this case, in order to achieve high spatial resolution, the laser beam is focused into a $\sim 1.5\,\mu\text{m}$ diameter spot, and the laser intensity is modulated at a frequency of about 0.1 MHz to limit the lateral thermal spread to a few μm [98, 99] (see Sec. 6.1.5). This approach is applicable as long as the thermal cycling does not lead to irreversible processes. The present work deals precisely with the metastable process of persistent photoconductivity, and therefore, a lock-in technique based on laser modulation is not a viable option.

As we pointed out above, the photoexcitation is not only able to reveal the standing wave of the fundamental frequency, but also clearly depicts the double anti-node of the current density in the resonator at the second harmonic. Further compelling evidence demonstrating the correlation between the photoexcitation map and the actual current density in the resonator can be obtained through simulations (Sonnet software) of the sample's response at both the fundamental frequency and the second harmonic. The resonance frequencies obtained for fundamental and second harmonics are 5.05 and 10.02 GHz, without considering the kinetic inductance. The resulting current density maps are shown in Fig. 6.3 (g,h). The extracted current density profiles along the resonator are shown in Fig. 6.3 (i,j) as a continuous line for both the fundamental and the harmonic, and they are compared with the experimentally determined δf_{PE} data (solid circles). The good agreement between the numerically determined current density and the experimentally measured photoexcitation sensitivity, for both fundamental and second harmonic resonances, validates the

interpretation discussed above.

6.1.4 Intermediate conclusion

In this first part of Chapter, we demonstrated the potential to locally modify the properties of oxygen-depleted YBCO resonators through photodoping. Several key findings merit emphasis. First, unlike traditional bolometric excitation, which is fast, volatile, and highly diffusive, the photodoping process is slow, stable, and non-diffusive. Despite this, it enabled us to map the current density within the resonators, as evidenced by imaging standing waves at the fundamental and second harmonic frequencies. Second, our approach can identify and highlight regions that are more susceptible to changes in the resonator's response. Third, the photodoped areas can be detected via changes in reflectivity using a photodoping-inactive wavelength or low-photon doses. While this technique is presented as a proof of concept, with clear opportunities for refinement, particularly in reducing the laser spot size and increasing scanning speed, it demonstrates great promise for direct and erasable laser-based scribing of resonators with tailored precision. Additionally, the proposed technique can be applied to other cuprates exhibiting photoexcitation, such as oxygen-deficient $\text{GdBa}_2\text{Cu}_3\text{O}_x$ [180] and likely $\text{Bi}_2\text{Sr}_2\text{Ca}_2\text{Cu}_3\text{O}_{10+x}$ [181], and could be extended to study persistent photoconductivity in strontium titanate [182] or lanthanum aluminate [183].

6.1.5 Annex

Magnetic field dependence of the resonance frequency

In order to probe the magnetic field dependence of the resonator response, the resonance frequency as a function of out-of-plane (OOP) magnetic field $\mu_0 H_{OOP}$ for zero-field cooling (ZFC) and field cooling (FC) has been measured. The results for the optimally doped YBCO at $T = 70\text{ K}$ are shown in Fig. 6.4. The response is similar to that presented in Ref. [172]. According to these authors, during

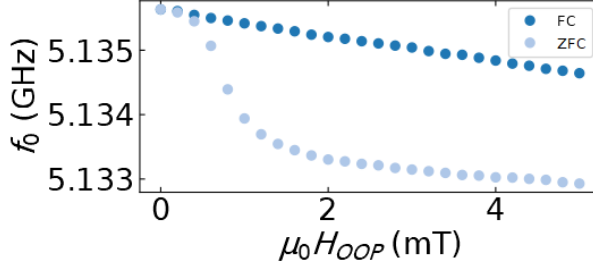


Figure 6.4: Dependence of the resonance frequency on the OOP magnetic history of the sample. FC shows a monotonic decrease as fluxons homogeneously populate the sample, whereas ZFC exhibits a sharp drop around 1 mT due to screening currents and increased magnetic field at the edges.

FC measurements, where magnetic field gradients are absent, the resonance characteristics are determined only by the vortex dynamics, exhibiting a nearly linear decrease with field. In contrast, in ZFC measurements, where highly inhomogeneous field distributions exist, the effect of screening current becomes relevant and gives rise to a more complex field dependence.

Transitory to stationary temperature evolution due to laser irradiation

When illuminating the sample with the laser, a non-uniform spatial temperature distribution is developed (see Sec. 6.1.5). In addition, this non-uniform distribution evolves over time, as it requires a finite amount of time for the heat diffusion to reach a steady state. During the procedure to obtain the maps presented in Fig. 6.3(c-f), it is essential to ensure that the sample has reached the steady state before measuring the resonance frequency. To that end, the time evolution of the resonance frequency needs to be characterized. This is achieved by measuring the transmission coefficient S_{21} as a function of time with a 4 ms time step at fixed frequency. This procedure is repeated for

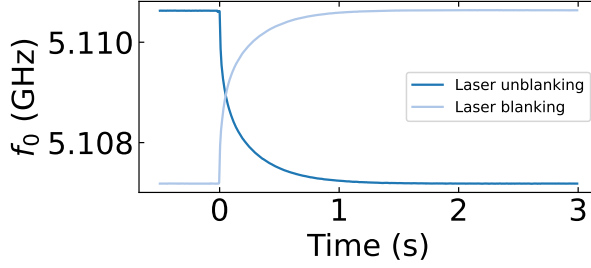


Figure 6.5: Time dependence of the resonance frequency when blanking/unblanking the laser spot parked at the center of the resonator. In both cases, the resonance frequency saturates in approximately 1 s.

equally spaced frequencies around the resonance frequency to obtain a matrix $S_{21}(t, f)$. By transposing the data, it is possible to reconstruct $S_{21}(f)$ as a function of time. A similar procedure was performed to measure the cooling-down process when the laser is blanked after reaching the stationary state. The results are summarized in Fig. 6.5.

Both during the heating (unblanking) and the cooling (blanking) processes, the resonance frequency saturates to its final value in approximately 1 s with a characteristic time constant of 111 ms, indicating that the sample has reached its thermal equilibrium.

Change in reflectivity

In addition to reducing the normal state resistance and increasing the critical temperature of oxygen-deficient YBCO, photodoping also increases its reflectivity [168]. To demonstrate this, the laser was parked at a fixed position in the ground plane, and images were taken over time. The mean value of the light intensity is measured in the spot area, and the relative change, equivalent to the change in reflectivity, is plotted in Fig. 6.6. The change in reflectivity tends to saturate over hours, similarly to T_c and ρ .

In a similar manner to how we previously detected minute shifts

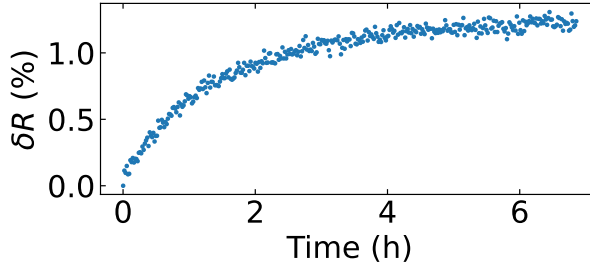


Figure 6.6: Relative change of reflectivity δR as a function of time acquired by parking the laser and tracking the change in gray value in the center of the laser spot. The light intensity is recorded using a CCD camera.

in the resonance frequency induced by laser irradiation on timescales much shorter than the characteristic photoexcitation time, we can also monitor changes in the sample's reflectivity during the laser scan. While the sample is continuously exposed to low-power red light, spatial variations in reflectivity are tracked in real time. To achieve this, images were acquired with the laser completely blanked (see Fig. 6.3(a)) at successive stages of the scan. These images were compiled into an animation (available in the Supplementary Materials of [47]), revealing a clear contrast between regions already scanned by the laser, where the reflectivity is noticeably enhanced, and regions yet to be irradiated. The average reflectivity of a set of regions along the slow scanning direction shows a jump in the reflectivity after the laser passes.

Local temperature rise due to laser heating

The change in resonance frequency due to laser heating, as shown in Fig. 6.5, could in principle be used to estimate the local increase in temperature. However, this is a complex problem requiring multiphysics modeling. Instead, to estimate the thermal spread during the laser

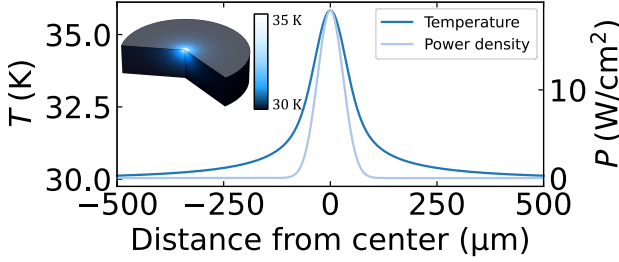


Figure 6.7: Numerical simulation of the heat distribution induced by a nonuniform Gaussian beam heating modeling the laser spot. The heat propagates laterally to distances much larger than the spread of the laser. The inset shows a 3D view of the heat distribution in the thickness of the simulated bi-layer.

scanning process, the numerical resolution of the stationary heat transport equation $\nabla \cdot (\kappa \nabla T) = -Q$ via finite element COMSOL solver, has been implemented. The considered parameters are sample thickness $d = 100$ nm, thermal conductivity of the YBCO film $\kappa_{YBCO} = 0.258$ Wm⁻¹K⁻¹ [184], substrate thickness $d_{LSAT} = 500$ μm, thermal conductivity of the substrate $\kappa_{LSAT} = 1.2$ Wm⁻¹K⁻¹ [185], heat transfer resistance between film and substrate $h = 1.5 \cdot 10^{-7}$ m²K/W [186, 187], Gaussian laser spot diameter $D = 60$ μm, and light power 10 mW from which 86% is absorbed in the YBCO layer [178]. A Dirichlet boundary condition corresponding to $T = 30$ K is imposed at the bottom of the substrate. The resulting temperature distribution based on these parameters is shown in Fig. 6.7.

The modeling indicates that the local temperature rises by approximately 5 K at the center of the laser spot, and that the thermal spread, as characterized by the full width at half-maximum, increases to approximately 120 μm in diameter. However, the sharp Gaussian profile of the heat source turns into a much broader thermal response, such that the sample is heated by more than 0.5 K over a region 500 μm in diameter, which is approximately one third of the resonator

surface. Furthermore, when the laser is parked in the gap of the resonator, less light is absorbed by the LSAT, i.e. approximately one fourth compared to the YBCO film. This ultimately produces approximately 4 times less heat and the maximum temperature rises by 1.2 K.

Validation of the phenomenological model

Fig. 6.8 shows a visual representation of the kernels and convolutions described in the phenomenological model above.

Fig. 6.9 compares the profiles from experimental frequency shifts with those from the phenomenological model described above. Because of its highly local nature, both the measured and modeled photoexcitation-induced frequency shifts shown in panels (a,b) reflect the local current density. They display a single anti-node sine function for the fundamental resonance frequency (panel (a)) and a double anti-node for the second harmonic (panel (b)). The thermally-induced frequency shifts, shown in panels (c,d), do not accurately reproduce the current density due to heat's diffusive nature. This results in maximum shifts at the device's center for the fundamental resonance frequency (panel (c)). Because of the meandering structure, the profile inherently oscillates as it enters and leaves the device center. In the second harmonic case (panel (d)), the maxima of the thermally-induced shifts occur where the current density peaks, i.e., at 2 mm and 6 mm. Despite minimal current at the resonator's center, a relatively high frequency shift is still detected when the laser is parked there. This occurs because the heat generated at the center node diffuses toward the two anti-nodes.

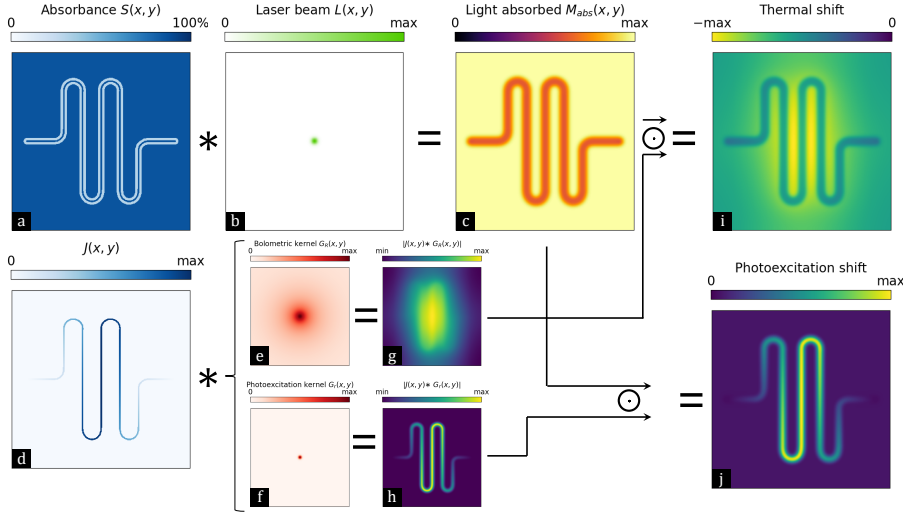


Figure 6.8: Convolution-based modeling of laser-induced frequency shifts. The absorbed light map $M_{abs}(x,y)$ (c) is obtained by convolving ($*$) the local absorbance of the sample $S(x,y)$ (a) (20% for LSAT and 86% for YBCO) with the Gaussian laser profile $L(x,y)$ (b). The spatial sensitivity of the resonator, which quantifies how strongly absorption at a given point affects the resonance frequency, is calculated by convolving the current density distribution $J(x,y)$ (d) at resonance, with spatial kernels representative of the underlying physical mechanisms. For the bolometric contribution, a broad thermal response function produces $G_R(x,y)$ (e), whereas a Gaussian profile matching the laser spot size produces $G_r(x,y)$ (f) for direct photoexcitation. Finally, the absorbed-light map (c) and the sensitivity maps (g,h) are multiplied element-wise ($\odot$) to obtain the expected spatial distribution of the frequency shift (i,j), which occurs only where absorbed light overlaps with resonator sensitivity.

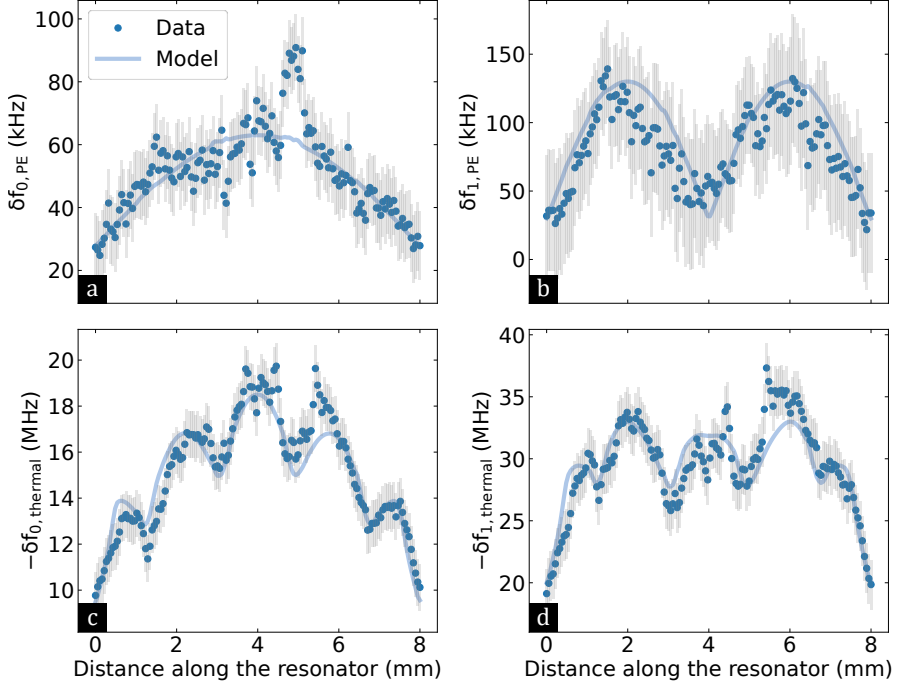


Figure 6.9: (a-b) Photoexcitation-induced and thermally-induced (c,d) frequency shifts measured (points) and modeled (continuous lines) extracted along a path in the center of the main meander of the resonator. The photoexcitation-induced frequency shift faithfully reproduces the current density in the device. Since the model is purely qualitative, the extracted profiles have been offset to correlate with the experimental data.

6.2 Time-resolved SLM

6.2.1 Introduction

Unlike wide-field imaging, scanning techniques suffer from inherently low speed acquisition, typically in the range of 10^3 px/s, which translates into a few seconds for a 512×512 px² image. Therefore, they are well-suited to investigating the quasi-static properties of the sample but offer limited access to short-time-scale processes. Achieving high spatial resolution while capturing dynamic effects still remains a challenging task.

A possible way to overcome this limitation is the so-called lazy fisherman method [188–191], in which the probe remains stationary at a single point while the system is driven by a periodic signal (electric, magnetic, elastic, etc.). Using the AC reference signal to set the time counter and assuming that the system response is also periodic, one can combine the periodic signal picked up at each pixel to construct a two-dimensional map revealing the oscillatory response of the system over time. If the excitation is not periodic, the collected signal represents a time-average, as illustrated for instance in Ref. [192–194].

An important extension of earlier approaches for systems driven by periodic excitations was introduced by Borgani *et al.* [195] to achieve sub-millisecond time-resolved conductive atomic force microscopy measurements. In this method, the sample is excited by an amplitude-modulated, square-wave–focused laser beam at a fundamental frequency ω_0 , while a large number of harmonic components of the photoresponse at integer multiples $n\omega_0$ are simultaneously acquired using lock-in detection. The frequency-domain data thus obtained can then be transformed into the corresponding time-domain response through an inverse Fourier transform. A somewhat similar approach has been used to reconstruct magnetization loops from magnetic susceptibility harmonic data [196, 197].

In this section, we apply this concept using hundreds of harmonics to reconstruct thermal-susceptibility images of a superconducting

YBCO resonator with a few- μ s time resolution, demonstrating a powerful and broadly applicable platform for resolving rapid thermal phenomena in condensed-matter systems. Both the scanning laser microscope used to acquire the data presented below and the multi-harmonic lock-in reconstruction scheme have already been described in detail in Chapter 4; only their application to the superconducting resonator is presented here.

We briefly recall the principle of the reconstruction. The photoreponse is chopped periodically at a fundamental frequency f , and a set of harmonics nf is collected simultaneously with lock-in amplifiers, each harmonic returning the complex Fourier coefficient C_n of the pixel signal. Applying an inverse Fourier transform pixel-by-pixel then yields the time evolution of the local thermal susceptibility within one chopper period $1/f$. For N collected harmonics the reconstruction provides frames spaced in time by $\delta t \propto 1/(2Nf)$, so that the achievable temporal resolution is set by the highest measurable harmonic rather than by the acquisition speed. The method, its validation on a square-wave photodiode signal, and the thermal-wave picture underlying the reconstructed dynamics are detailed in Chapter 4 (Eq. (4.8) and Figs. 4.5 and 4.4). As emphasized there, the reconstructed frames map the macroscopic dynamic susceptibility to the local stimulus and should not be read as a direct thermograph of the local heat relaxation.

6.2.2 Case study: superconducting resonator

The time-resolved reconstruction technique was implemented on the optimally-oxygenated $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ half-wavelength superconducting resonator described in Sec. 6.1. The specimen was mounted within a Montana Instruments optical cryostat. One port of the resonator was coupled to a Vector Network Analyzer (VNA) operating in continuous wave (CW) mode, while the output port was connected to a 15 dB RF amplifier followed by a Krytar 203B RF diode. The diode rectifies the RF signal into a DC voltage proportional to the transmitted power.

This voltage was subsequently processed by a lock-in amplifier phase-locked to the reference frequency of the SLM optical chopper. To characterize the device, the scattering parameter $S_{21}(f)$ was initially recorded to identify the fundamental resonance frequency. The inset of Fig. 6.10(a) displays the resonance peaks obtained at $T = 30$ K (black) and $T = 32$ K (yellow). Note that the resonance frequency shifts to lower frequencies with increasing temperature, a behavior attributed to the enhancement of the material kinetic inductance, as discussed in Sec. 6.1. This redshift leads to a measurable change in DC output δ of the RF diode, which is used as photoresponse. This preliminary characterization facilitates the selection of an optimal probe frequency for the LIA. During SLM operation, the modulated laser beam induces a periodic shift in the resonance peak (as illustrated in the inset) thereby modulating the rectified DC output voltage at the chopper frequency. In the perturbative regime, wherein the laser irradiation does not induce significant bolometric heating of the device, the modulation of the transmitted power remains directly proportional to the shift in the resonance frequency.

Time-resolved response

As previously detailed, SLM images were acquired using a chopping frequency of 180 Hz. To achieve a comprehensive time-resolved reconstruction, 300 harmonics were recorded simultaneously by employing three LIAs in parallel, with each instrument dedicated to a specific harmonic range: 1–100, 101–200, and 201–300. The high input impedance of the SR830 (10 M Ω) facilitates the parallel coupling of multiple instruments without inducing signal distortion or loading effects. Consequently, the total acquisition time is reduced by a factor of approximately $N = 3$, where N represents the number of instruments utilized in parallel. To ensure precise temporal synchronization between the reconstructed frequency response and the optical excitation, the signal from the photodiode was monitored using a fourth LIA restricted to odd harmonics. As previously demonstrated in Fig. 4.5,

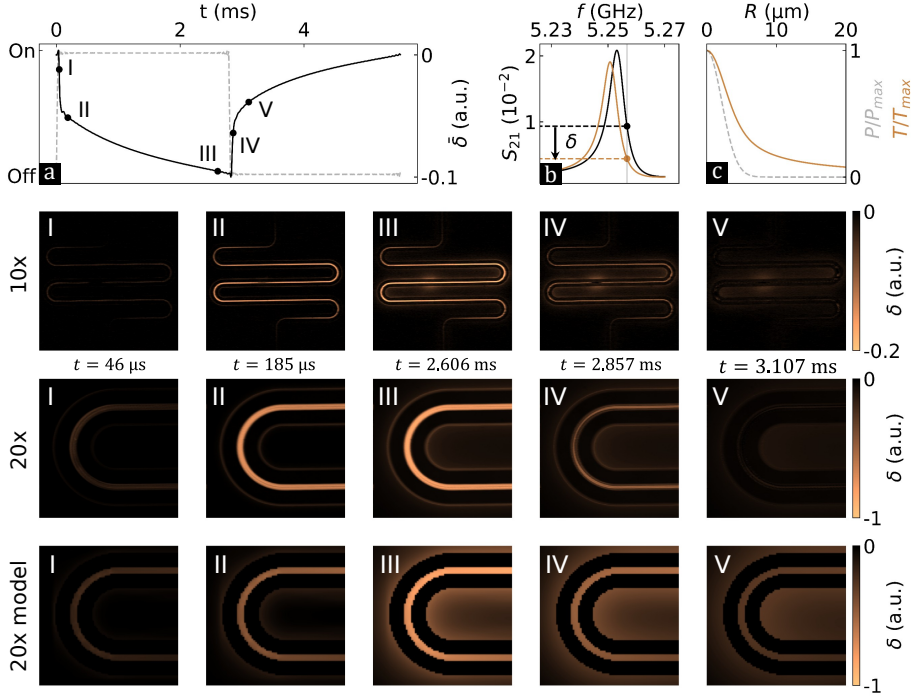


Figure 6.10: (a) Time-resolved reconstruction of the average response of the superconducting resonator over one chopper period. The signal captured by the photodiode, corresponding to the on/off state of the laser, is plotted in a dashed line. (b) Graphical representation of the probed photoresponse δ , corresponding to the change in RF power transmitted through the resonator at fixed frequency and which depends on the temperature distribution in the sample. The black (yellow) curve illustrates the case where the laser is off (on). (c) Optical (dashed gray) and thermal (yellow) spot profiles. The thermal profile is obtained from simulation and is used as the kernel for the convolution model of the bottom row. Snapshots at five different instants, labeled I-V and indicated in panel (a), are shown in the bottom three rows. The first row corresponds to images of the photoresponse of the entire sample using a $10\times$ objective, whereas the second row shows images collected using a $20\times$ objective, zooming on the bottom left half-turn of the meander. The bottom row shows the result of the time-dependent convolution model developed in the text.

collecting even harmonics is not necessary as they are weak compared to the odd ones. This configuration provides sufficient resolution to accurately determine the onset of the laser-induced transition.

The spatially averaged microwave response of the resonator is presented in Fig. 6.10. Upon laser excitation, an initial rapid response is observed within 100 μs . The relative shift δ is negative, as the continuous-wave probing frequency is situated on the high-frequency slope of the resonance; consequently, the thermally induced redshift of the resonance peak results in a decrease in transmission (see Fig. 6.10(b)). Following this initial onset, δ continues to decrease at a significantly lower rate.

To elucidate this two-step behavior, we examine five temporal snapshots showing $\delta(x, y)$ captured during the periodic stimulus. These images represent time-resolved maps of the local photothermal susceptibility. Immediately following the onset of irradiation (I), the response is localized strictly to the meander. Since the central conductor carries the highest current density (proportional to the quality factor), it exhibits the highest sensitivity and yields the fastest response time. This characteristic timescale likely scales as a^2/D , where a denotes the beam waist and D is the thermal diffusivity.

After 185 μs (II), the susceptibility of the meander reaches saturation, while the edges of the ground plane, characterized by lower current densities, begin to emerge. Indeed, the edge currents in the ground plane, next to the excitation line or the resonators, are about a factor $1 + 2g/w \sim 2$ smaller than in the central line [98], where g is the gap between central track and ground plane, and w is the width of the central track. By the conclusion of the heating phase (III), the central conductor remains saturated, whereas the ground plane exhibits a pronounced susceptibility. This suggests that over these extended timescales, heat deposited in the middle of the ground plane diffuses through the substrate to the central conductor, resulting in a delayed contribution to the reconstructed signal.

This interpretation is further supported by the dynamics observed immediately following laser extinction (IV): the susceptibility of the

central conductor decays rapidly, while the signal from the ground plane persists. This behavior demonstrates a clear phase delay governed by the finite thermal diffusivity of the system. Finally, the cycle concludes as heat dissipates through both the film and the underlying substrate.

Modeling

To substantiate these claims, a modeling approach analogous to that described in Sec. 6.1 was implemented. This model relies on the principle that the resonance frequency of the resonator is determined by the total capacitance and inductance of the device. Because the inductance contains a kinetic contribution from the superconducting condensate, its value depends directly on the local temperature. Furthermore, the magnitude of the resonance frequency shift is proportional to the current density circulating through the regions of varying kinetic inductance. Consequently, local temperature fluctuations in areas supporting negligible current density yield no observable effect on the global response, and conversely. This rationale motivates the introduction of a convolution product between the local current density and the spatial temperature profile induced by the laser. To implement this, COMSOL Multiphysics was employed to simulate the time-dependent temperature profile generated by the laser illumination. The numerical simulation modeled a 500 μm -thick LSAT slab (with parameters $\kappa = 1.5 \text{ W/mK}$, $C_v = 200 \text{ J/kgK}$, and $\rho = 6740 \text{ kg/m}^3$) thermalized to 30 K at the bottom surface. Those parameters were optimized to match the experimental response of the sample. Unfortunately, little thermal properties measurement for this particular substrate were found in the literature [185].

Focus on the delayed response

To conclude the analysis of the time-resolved response, we focus on the slight delay that develops between the excitation and the response when the laser is positioned far from the sensitive regions.

As recalled in Sec. 4.2.2, the harmonics of the square-wave input all propagate at different speeds and are damped exponentially more strongly at shorter wavelengths (*i.e.*, higher harmonic excitations). This progressively distorts the response, which evolves from a square wave into a sawtooth and eventually retains only the fundamental sine component. Moreover, the differing phase velocities introduce a delay between the temperature extrema close to the source and those far from it.

Fig. 6.11 reproduces a snapshot of the resonator response introduced above, while its normalized local response, probed as a function of time in several regions, is shown in Fig. 6.11(b). The light-blue region is measured within the resonator meander, where the response most closely follows the excitation. The farther the probing point lies from the sensitive regions, the more the waveform is distorted, gradually adopting the shape of the intermediate regime presented in Fig. 4.4.

We next turn to the delay between the switching on of the laser and the maximum of the response. Fig. 6.11(c) maps this delay, extracted for each pixel from the time-resolved response. This delay is obtained by fitting the response with an empirically chosen periodic saturating-exponential function with two characteristic timescales: a short one accounting for the rapid variations, and a longer one, of the order of the total period, accounting for the global response. When the laser is parked at the center of the meander, the delay is virtually zero, since the response is synchronized with the excitation. As the illumination point moves deeper into the ground plane, toward the center of the image, a delay builds up, visible in the dark-blue curves of Fig. 6.11(b), amounting to a few percent of the total cycle time. Cross-sectional profiles of this delay as a function of distance r are shown in Fig. 6.11(d), where $r = 0$ denotes the point closest to the meander half-turn. As expected, the delay increases as one moves away from the current-carrying parts, and two distinct trends emerge. The red cross-section shows a delay that first increases and then reaches a plateau. This can be understood as follows: for values

of r up to the meander radius, the distance to the nearest current-carrying part grows, so the delay grows accordingly. Beyond this region, the area most strongly affected by the laser is no longer the half-turn but the straight portions of the resonator, whose distance to the cross-section line remains constant as r increases; the delay therefore saturates.

Conversely, the orange cut line measures the delay radially outward from the center of the resonator. In this case, as r increases, the entire current distribution of the resonator recedes from the probing point. The delay consequently increases linearly, since the current-carrying region reached first by the heat is the one located at $r = 0$. Fitting this delay with a linear law yields a propagation speed of $307 \pm 4 \mu\text{m/ms}$. As noted in Sec. 4.2.2, although heat itself does not propagate like a wave, its modulation does. The speed at which the fundamental propagates, given by $v = \sqrt{2D\omega}$, is, to first approximation, the one that gives rise to the measured delay. The higher harmonics are neglected here, as they are damped more strongly and thus contribute less to the delay.

From the above expression for the speed of the fundamental harmonic, a diffusivity of $41.7 \pm 1.2 \text{ mm}^2/\text{s}$ is found. This value differs markedly from the one used in the COMSOL simulations presented above, $D = \kappa/\rho C_v \approx 1.1 \text{ mm}^2/\text{s}$, being about a factor of 40 larger. Such a discrepancy is nonetheless unsurprising, as the present problem is considerably more complex than the model used to describe it. As noted earlier, all interfacial thermal effects were neglected in the modeling, and the parameters were optimized to reproduce the experimental response as closely as possible. As a result, the values obtained should be regarded as effective quantities rather than as absolute, metrological measurements.

Care must be exercised when interpreting these results. The probed quantity is the global shift of the resonance frequency of the resonator, obtained while illuminating a spot orders of magnitude smaller than the device itself. This minute spot generates a thermal profile that extends well beyond its own size and takes time to diffuse.

Over this diffusion length, several regions of the resonator are heated, some of which may carry a high current density. It is, however, difficult to disentangle the contribution of a strongly current-carrying region heated only slightly from that of a moderately current-carrying region heated substantially, as both may produce the same frequency shift. In other words, the notion of *sensitive region* introduced earlier is more of a conceptual construct than a rigorously definable quantity.

6.2.3 Conclusion

In summary, we propose an experimental approach to obtain time-dependent scanning laser microscopy images by periodically illuminating each pixel, collecting multiple harmonics with a tandem of lock-in amplifiers, and then applying an inverse Fourier transform. We first validate the technique by reconstructing the square-wave photoresponse from an amplitude-modulated laser source. Subsequently, we apply this method to image a standing wave at the fundamental resonance frequency of a YBCO resonator as a function of time during laser irradiation. Regions with higher current density are more sensitive to local heating and provide a sizable photoresponse signal. The contrast variation of the images over time reflects the spreading of the thermal spot due to thermal diffusion, whose regimes are described by the thermal-wave analysis of Chapter 4. In principle, this technique permits estimation of the thermal parameters governing heat diffusion while revealing the more thermally susceptible regions of the sample.

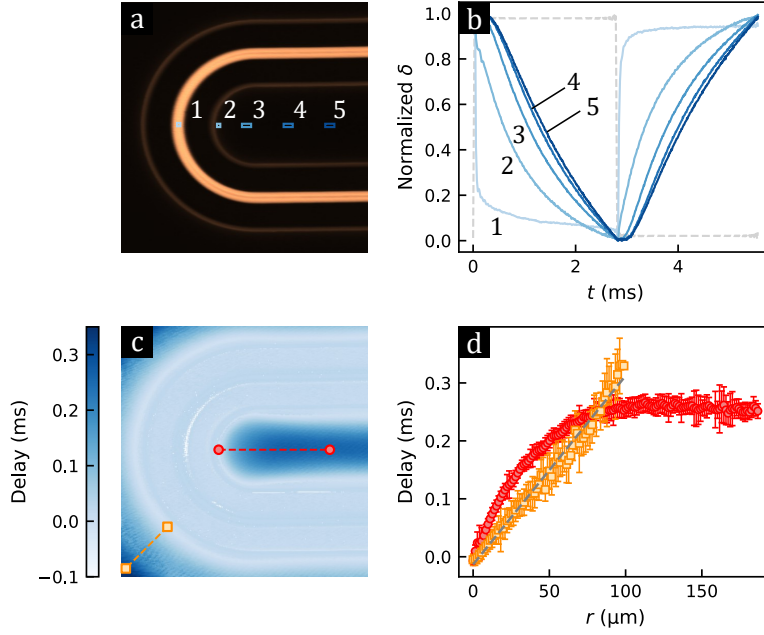


Figure 6.11: (a) Response of the resonator $55\ \mu\text{s}$ after illumination. (b) Time-resolved local response of the resonator. Each probed area is marked in (a) with a rectangle. The response is normalized, and the reconstructed photodiode signal is shown as a dashed gray line. The dark-blue curves, deep in the ground plane, exhibit a pronounced delay between the switching of the laser and the maximal response. (c) Delay extracted pixel by pixel from the temporal response; the darker a region, the longer the time required for the heat to propagate from that pixel and affect the global resonance frequency. (d) Profiles extracted from (c) along a line in the central ground plane (red) and along a line in the ground plane extending away from the resonator (orange). The red curve tends to saturate because the probed points remain within the resonator structure, between the parallel arms of the meander, whereas along the orange line the delay increases linearly with distance, as the probed point moves farther from all current-carrying regions. A propagation speed is extracted from the linear fit of the orange data, shown in gray.

Chapter 7

Steering Magnetic Flux with Laser-Induced Periodic Surface Structures

The preceding chapters have shown that penetrating magnetic flux is, almost invariably, detrimental to the performance of a superconducting device: in the NbTiN resonators of Chapter 5, vortices and thermomagnetic flux avalanches affect the resonance frequency and degrade the quality factor. The natural response, pursued there, is to *keep the magnetic flux out*. This chapter takes the complementary stance: rather than merely excluding flux, we wonder whether its penetration can be actively *steered*, so that flux is guided along, or diverted away from, chosen regions of a superconducting film. The tool used to achieve this is a purely topographic one, namely laser-induced periodic surface structures (LIPSS), and the physics is read out with the same magneto-optical imaging (MOI) technique developed in Chapter 3.

Custom-shaped magnetic flux guiding channels are fabricated on superconducting Nb thin films by laser nanopatterning of their surface. Preferential pathways are defined by suitably combining imprinted anisotropic pinning domains through LIPSS. Generated by

the selective energy deposition of femtosecond UV laser pulses, quasi-parallel ripple structures form under optimized irradiation conditions; on average, each domain is composed of grooves with a lateral period of 260–270 nm and a depth of about 80 nm. By combining scanning and transmission electron microscopy, magneto-optical imaging, and conductive atomic force microscopy, we conclude that the boundaries of the LIPSS-covered domains play a prominent role in the magnetic flux diversion process within the film. This is confirmed by dedicated modeling of the flux dynamics, combined with the inversion of the magneto-optical signal. The created metasurfaces enable control of the flux penetration process at the microscale.

A note on the origin of this chapter. This chapter is adapted from E. Martínez & N. Lejeune *et al.*, “Laser engineered architectures for magnetic flux manipulation on superconducting Nb thin films” [77], and builds directly on the laser-nanostructuring study of Badía-Majós *et al.* [198], from which the anisotropic critical-current picture is inherited. The magneto-optical imaging technique and the Biot–Savart current-density inversion used throughout are developed in Chapter 3.

Personal contribution. The Nb thin films were deposited in Liège and subsequently laser-nanopatterned at the University of Zaragoza, where the critical-state modeling of the anisotropic flux dynamics was also developed. I carried out the magneto-optical imaging of the architectures presented here, as well as the Biot–Savart inversion of the magneto-optical signal into current-density maps on which the interpretation rests. The analysis of the flux penetration and diversion mechanisms was led by myself, with the involvement of all co-authors.

7.1 Introduction

The manipulation of electromagnetic (EM) fields plays a prominent role in the development of cutting-edge technologies, in particular within the field of microelectronics. Of special mention are the superconducting circuits for quantum computation and magnetic sensing [199, 200], the nano-magnet based logical circuits [201, 202], and the so-called fluxonic devices [203–206], that may host conceptual breakthroughs in computer engineering. Switching the system from one state to another and reaching enough sensitivity require control of the applied magnetic fields in the submicron scale, a target that may be achieved by superconducting nanoarchitectures. Although all these applications take advantage of many properties related to the superconducting state (zero dissipation, fast switching velocities, magnetic flux quantization), one must also face a variety of shortcomings. In particular, let us mention that the nucleation of quantized magnetic flux units (vortices) [53] may become highly detrimental when they penetrate in the form of catastrophic flux avalanches. This has been observed in the superconducting radio-frequency resonators [28, 42] that are the subject of Chapter 5, where such avalanches were shown to trigger abrupt jumps in the resonance frequency. Other issues to be accounted for relate to geometrical effects, highly relevant due to the quasi-2D nature of the integrated technologies. Thus, finite size effects with the concomitant demagnetizing fields or the rapid spread-out of field with distance require dedicated consideration if either field concentration or shielding in a targeted area is to be achieved.

Focused on superconducting technologies, the idea of guiding magnetic fields by using artificial structures acting as flux rails comes from aged contributions [207–210] showing that, in thin samples, perpendicular magnetic flux vortices tend to be guided along microstructural features such as grooves in grated structures. Characteristic grating (line separation and depth) must be in the submicron scale, since the basic underlying physical principle is the reduction of vortex-line energy, being their typical dimensions in the range of tens or

few hundreds of nanometers. More recently, the advent of powerful nanofabrication techniques has allowed a further development and deployment of this concept. For example, uniaxial nanogroove structures have been used to demonstrate the transition from isotropic to anisotropic flux penetration regimes by tuning up the working temperature [211] and a DC-tunable microwave low-pass filter has been achieved by the construction of asymmetric nanogroove morphologies [205]. Taking a step further, here we put the focus on the extension of the magnetic flux guiding process along quasi-2D channel architectures, i.e. on the fabrication of pathways enabling the manipulation of the magnetic flux piercing the superconducting film, either by concentrating or diverting it from a specific area of the device. Superconducting Nb thin films have been selected for this study for their favorable properties (highest T_c of elemental superconducting materials, type-II, widely used and grown worldwide) and technological relevance. Nb finds applications in planar resonators [28], particle accelerators [212], etc., for which relevant physical effects could be affected by surface nano-structuring. The application of surface tailoring technologies to materials with higher complexity, in particular to the high-temperature superconductors studied in Chapter 6, is still a challenge to be addressed in future research.

Our proposal relies on a one-step processing approach that functionalizes the material by modulation of its surface topography in the submicron scale. Taking advantage of recent progress in laser technologies for material processing, one may induce corrugation of surfaces with control over: (i) the lateral periodicity, (ii) the depth, and (iii) the local orientation of the topographical grooves, all these with minimal alteration of the material itself. In fact, direct laser processing enables a versatile matter manipulation via irradiation with ultrashort (femtosecond) controllable laser pulses. Remarkably, owing to the intrinsic thermalization processes related to the interaction of matter and ultrafast laser pulses, one gets rid of unwanted heating effects [213]. Therefore, material structure preserving modifications, such as the good match between film and substrate, are allowed. In

particular, this processing method has been used to produce quasi-parallel surface patterns, the so-called Laser-Induced Periodic Surface Structures, LIPSS. Often termed simply “ripples”, LIPSS are a universal phenomenon that manifests during the laser processing of solids upon irradiation with coherent laser radiation [214–216]. Hitherto, it has been applied to different materials and for different functionalities such as iridescence of surfaces, wetting control (hydrophobic-philic surfaces), medical cell growth or tribological uses. Details on these applications may be found in Refs. [217, 218]. Phenomenologically, quasi-periodic arrangements of rippled lines or other surface features with spatial periods Λ of the order of the irradiation wavelength λ or even far below are formed in “self-ordered” ways after multiple-irradiation events in the focus of the pulsed laser beam. LIPSS arise from coherent optical scattering and (intra-pulse) interference effects at nanoscale surface defects along with subsequent positive (inter-pulse) feedback [217, 219]. Even if the seeding of LIPSS always arises from electromagnetic effects, subsequent matter-reorganization (e.g. via thermal melting and hydrodynamic effects) may additionally affect the final surface topographies [216, 220–224]. LIPSS, being already present in the laser irradiated spot, can be “coherently linked” together [225] via laser beam scanning approaches for covering the desired macroscopic surface areas.

For the metal niobium, different types of LIPSS were reported in the literature. Most prominently, so-called low spatial frequency LIPSS (LSFL) with periods $\lambda/2 \leq \Lambda \leq \lambda$ form in the direction perpendicular to the linear laser beam polarization for sub-ps laser irradiation [226–229]. It is generally accepted that the LSFL originate from coherent scattering and interference effects of the incident laser radiation at the microscopically rough surface, eventually leading to a spatially modulated absorption of the optical radiation. For suitable optical properties of the irradiated material, this may include the excitation of Surface Plasmon Polaritons (SPPs) [216]. Moreover, Pan *et al.* reported the formation of high spatial frequency LIPSS (HSFL) with 290–350 nm sub-wavelength spatial periods ($\Lambda \ll \lambda$)

being parallel to the laser beam polarization, upon irradiation with near-infrared laser radiation ($\lambda = 800$ nm) using pulses with a duration of 100 fs [226].

Outstandingly, the pulsed-laser irradiation technique offers important advantages with respect to alternative methods such as heavy ion irradiation, nanolithography or chemical routes. In particular, it is eco-friendly, scalable and efficient, allowing large surfaces to be covered within a short time, and it has the possibility to tailor the pinning landscape at will and with high resolution. Furthermore, the presence of undesired chemical reactions, such as oxidation of the metal film, may be avoided by performing the process under a controlled inert atmosphere.

As a matter of fact, it is known that the morphological transformation of superconducting Nb induced by laser patterning gives way to conspicuous anisotropy in its electromagnetic response [198, 227]. In the case of thin films, LIPSS induce a noticeable anisotropy factor $J_{c\parallel}/J_{c\perp}$ in the ratio of the critical superconducting current densities flowing either parallel ($J_{\parallel}$) or perpendicular ($J_{\perp}$) to their direction [198]. Thus, the electric current density $\mathbf{J}$ flows preferentially along the LIPSS ridges. In addition, recalling that in type-II superconductors, flux penetration takes place in the form of incoming vortices (magnetic flux quantum units, i.e. Φ_0) boosted by the Lorentz-like force and restrained by the pinning barriers, defined by the topographical undulations themselves, one straightforwardly finds that flux preferentially penetrates along the LIPSS.

In this chapter we explore the magnetic flux guiding capacity induced by imprinting conjoined domains of surface LIPSS with appropriate orientations so as to control locally the critical current density. We will always refer to LSFL-type LIPSS for which the laser processing parameters were optimized. Dedicated simulation studies based on the Nb optical response confirm this working hypothesis (not reproduced here; see the supplementary materials of Ref. [77]). Nanocrystalline Nb thin films deposited on a silicon wafer, with initial thickness of 200 nm, were processed with an ultraviolet

pulsed laser (343 nm wavelength, 238 fs pulse duration). The flux channeling effects and their physical origin in these 2D structures are studied by magneto-optical imaging (MOI), together with the data inversion techniques introduced in Section 3.3.2 applied to the experimental profiles, and analyzed in terms of the critical state model (CS model) [230, 231] for the penetration of magnetic flux in type-II superconductors. Scanning and transmission electron microscopy (SEM, TEM) as well as topographic and conductive atomic force microscopy (AFM & cAFM) data are used in order to verify the surface topography, and to analyze the correspondence between topographical characteristics and magnetic flux profiles. Global superconducting properties of the laser processed films are characterized by SQUID magnetometry.

7.2 Materials and Methods

7.2.1 Fabrication

Film deposition and sample shaping

Thin films of Nb with a thickness of 200 nm were deposited onto Si/SiO₂ wafers, using radio-frequency magnetron sputtering. The Nb physical vapor deposition was conducted at a rate of 0.1 nm/s under an Ar pressure of 5 Torr, after pumping down the chamber to a base pressure of $4 \cdot 10^{-8}$ Torr.

Before laser patterning, $4 \times 4 \text{ mm}^2$ samples were laser-cut using a pulsed green laser (Light Conversion, Vilnius, Lithuania, model Carbide CB3-40W, $\lambda = 515 \text{ nm}$, $\tau_p = 249 \text{ fs}$) working at a pulse repetition frequency $f_{\text{rep}} = 10 \text{ kHz}$ and a laser scanning speed of 5 mm/s. Subsequent to the cutting process, the samples were ultrasound cleaned in isopropanol for 15 min.

Laser patterning

For the surface nanopatterning of the thin film, the above described laser with a wavelength of $\lambda = 343$ nm was used, along with a pulse duration $\tau_p = 238$ fs. At the working distance, the output laser beam exhibited a spatially Gaussian energy profile with an elliptical spot, where $1/e^2$ beam diameters were $2a \times 2b = 36 \times 60 \mu\text{m}^2$. An optimization process was performed in order to determine the laser parameters for sample processing, resulting in pulse energies in the range 3.40–3.98 μJ , laser scan speed of 125 mm/s, pulse repetition frequency $f_{\text{rep}} = 20$ kHz, and distance between scan lines of 8 μm , for the different samples under study, and in a controlled Ar atmosphere (under gas flow in an open chamber). The inert atmosphere was used in order to minimize the effect of oxidation processes that lead to perceptible reductions of the critical temperature T_c [198], not to mention the contribution to remove possible debris particles through the gas flow. The laser scan speed and the distance between scan lines were carefully selected for the above levels of pulse energies for obtaining a nearly uniform nanostructure distribution in the Nb thin film scanned area. A basic sketch of the laser-patterning process is included in Fig. 7.1. We notice that, in all cases, the beam scanning direction was parallel to the polarization of the beam. This relative orientation is not critical for the generation of LIPSS, although it may have effects on the overall morphology (specifically on their length) according to previous studies in other materials [232]. Then, in order to avoid data dispersion we were careful to keep the angle constant throughout the treatments.

In order to create the various patterns considered, subsequent to the calibration process, spatially resolved laser treatments were performed. Subsequently, the sample was rotated precisely by 90° using a goniometer and a home-made setup. Laser treatments were then conducted on the non-treated areas to produce LIPSS perpendicular to those generated previously, achieving the desired pattern. Finally, in order to obtain a centered pattern, as well as to avoid boundary

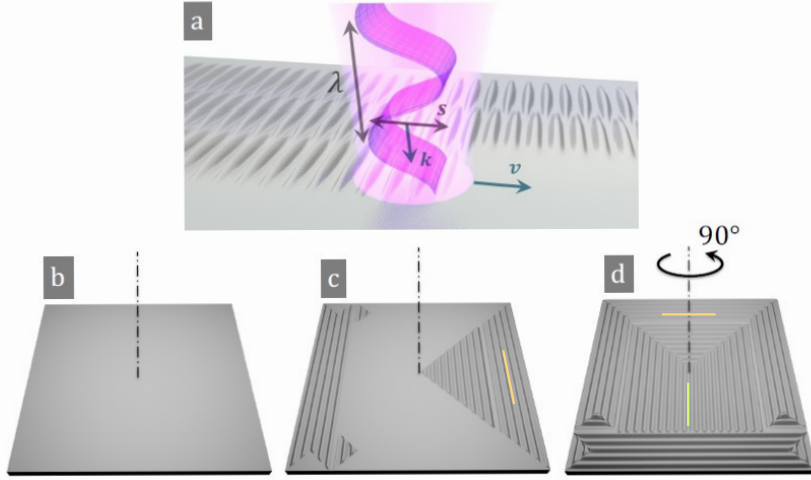


Figure 7.1: Sketch of the laser irradiation process. (a) The laser beam scans over the film surface, with its linear polarization aligned with the scanning direction. The corrugation of the surface emerges perpendicular to this direction. As sketched in (b–d), complex domain structures were created from the pristine sample (b) by using a two-step process according to pre-established configurations. This particular example corresponds to the sample that will be named after *top-hat*. Firstly, some selected areas were irradiated to generate LIPSS parallel to the orange segment (c). Secondly, the sample was rotated 90° and then (other) selected areas were irradiated to create LIPSS in the orthogonal direction, i.e. parallel to the green segment (d). Plain gray corresponds to non-irradiated areas while stripped areas are indicative of the laser-scanned parts.

effects due to imperfections in the cutting process, an additional laser treatment was applied aiming at the removal of Nb in excess, with inner dimensions $3 \times 3 \text{ mm}^2$ by applying an average power of 1.35 W to the original $4 \times 4 \text{ mm}^2$ platelet in a two-fold process. In brief, we end up with a perfect square-shaped film area with aligned LIPSS morphology, surrounded by a small inert frame of uncoated substrate.

7.2.2 Microstructural characterization

SEM and TEM analysis

Surface characterization was carried out using a MERLIN field-emission scanning electron microscope (FE-SEM) (Carl Zeiss in Jena, Germany), equipped with an energy dispersive X-Ray spectroscopy (EDS) system from Oxford Instruments (Abingdon, UK). The FE-SEM was operated at 5 kV and utilized secondary electron (SE), in-lens, and energy-selective backscattered (ESB) detectors. Cross-sectional TEM images were taken with a Tecnai F30 microscope of FEI (Lincoln, NE, USA). With this aim, several lamellas were prepared using a Focused Ion Beam (FIB) in a Dual Beam Helios 650 apparatus of FEI. Prior to the FIB-lamella extraction, samples were coated with C-Pt protective layers.

Atomic force microscopy (AFM and cAFM)

Atomic force microscopy measurements were performed with a scanning probe microscope (SPM) Ntegra Aura system of NT-MDT at room temperature in order to analyze the LIPSS topography. Quantitative local electrical resistance measurements were also performed at room temperature using the c-AFM mode of a CSI Instruments Nano-Observer AFM. The analysis of the AFM measurements as well as the 2D-Fast Fourier transforms from SEM images were carried out using Gwyddion software.

7.2.3 Magneto-optical imaging

The magnetic flux penetration along the laser-patterned areas was visualized by the magneto-optical imaging (MOI) technique, whose physical principle, cryogenic setup, and calibration are described in detail in Chapter 3. In brief, the local out-of-plane flux density $B_z(x, y)$ at the sample surface is mapped through the Faraday rotation imprinted on linearly polarized light traversing a Bi:YIG indicator placed in direct contact with the sample. The 3.5 μm -thick Bi:YIG active layer, epitaxially grown on a 450 μm -thick $\text{Gd}_3\text{Ga}_5\text{O}_{12}$ (GGG) substrate and backed by a 100 nm Al mirror, converts the field distribution into an intensity map recorded by a CCD camera. Since no radio-frequency excitation is involved here (in contrast to the resonator study of Chapter 5), the standard mirror-backed indicator is used without restriction.

7.2.4 Magnetic characterization

The complex AC susceptibility, $\chi_{ac}(T)$, with in-phase χ' and out-of-phase χ'' components, was measured at zero DC magnetic field and AC drive magnetic field (sine wave of amplitude $\mu_0 h_0 = 10 \mu\text{T}$, and frequency $f = 10 \text{ Hz}$) in a SQUID magnetometer (MPMS-5T from Quantum Design, San Diego, CA, USA), in order to determine the superconducting critical temperature T_c of the samples, using the onset of diamagnetism criterion. In pristine films we found $T_c = 8.9 \text{ K}$, while laser treated samples show a reduced value $T_c = 8.7 \text{ K}$.

All magnetic measurements were performed on the above described square samples with the magnetic field applied perpendicular to the surface and after zero-field cooling. For each measurement, a reset of the superconducting magnet was performed previous to sample cool down in order to eliminate the remnant magnetic field.

7.3 Critical-state modeling of the anisotropic flux landscape

Magnetic flux penetration in the laser patterned Nb samples has been simulated by a numerical extension of the CS model in type-II superconductors. In its original form [230, 231], the model postulates the *criticality* condition $J(x) = 0, \pm J_c$ for the circulating current density within the superconductor in 1D geometry (slabs and long cylinders). In its simplest form J_c , the so-called critical current density, is a material dependent parameter. Then, the actual magnetic field profile is established by integrating Ampère’s law, implementing boundary conditions on the sample’s surface and implicitly applying Faraday’s law for choosing the $+$, $-$ or 0 option. As a result, one obtains the actual flux penetration profile for an arbitrary route of the excitation field. In brief, this model is physically meaningful for type-II superconductors with strong flux pinning mechanisms. Then, the magnetic flux profile is basically determined by postulating the maximum gradient law ($J(x) = 0, \pm J_c$). This scenario is valid as long as the sample is well penetrated by “anchored” flux tubes and thermal agitation is negligible. In our experimental conditions (low temperatures and fields in the range of mT) both conditions are satisfied. In fact, for the case of flat samples, flux penetration occurs for fields well below H_{c1} , i.e. $H_{\text{ENTRY}} \approx H_{c1} \sqrt{0.36\zeta}$ [233], with ζ the platelet aspect ratio ($\simeq 7 \times 10^{-5}$ in our samples). On the other hand, the CS model will be used in upgraded form (2D model geometry for thin films in perpendicular field [234]), which is also justified by the small value of ζ . In this situation, not only the sign, but the local orientation of the current density is unknown at each point. In the quasi-steady regime, the electromagnetic problem is solved by taking benefit of the divergenceless property of the current density together with the negligible sample thickness d . As in the Biot–Savart inversion of Section 3.3.2, these conditions allow the problem to be posed in terms of a scalar function, here the *streamfunction* $\sigma(x, y)$ that relates to the surface current density (averaged over thickness)

by $\mathbf{K} = -\hat{z} \times \nabla \sigma$ [235], being

$$\mathbf{K}(x, y) = \int_{-d/2}^{d/2} \mathbf{J}(x, y) dz. \quad (7.1)$$

Finally, the application of Faraday's law together with the superconducting material law is accomplished by an optimization problem [234, 236]. One has to minimize the functional

$$\begin{aligned} \mathcal{F}[\sigma_{n+1}] \equiv & \frac{1}{2} \iint_{\Omega} \left[\frac{\nabla \sigma_{n+1}(\mathbf{r}) \cdot \nabla \sigma_{n+1}(\mathbf{r}')}{\|\mathbf{r} - \mathbf{r}'\|} - 2 \frac{\nabla \sigma_n(\mathbf{r}) \cdot \nabla \sigma_{n+1}(\mathbf{r}')}{\|\mathbf{r} - \mathbf{r}'\|} \right] d\mathbf{r} d\mathbf{r}' \\ & + 4\pi (H_z^{a,n+1} - H_z^{a,n}) \int_{\Omega} \sigma_{n+1}(\mathbf{r}) d\mathbf{r} \end{aligned} \quad (7.2)$$

with Ω the surface of the film, $H_z^{a,n}$ the z -component of the applied external field, and the subindices n and $n+1$ used to indicate two successive time layers. The superconducting material law (action of barriers for the flux dynamics that are codified by the critical value $K_c = J_c d$) enters through a restriction function that accompanies the minimization process and models the anisotropic and local character of the current density. In our case, related to the preferential flux slippage along given directions, it takes the form

$$\frac{K_{\parallel}^2}{\Gamma^2} + K_{\perp}^2 \leq K_c^2, \quad (7.3)$$

with $K_{\parallel}, K_{\perp}$ the current density components either parallel or perpendicular to the “easy direction” defined by the LIPSS at a given point. Minimization is performed step by step, adapting the response of the sample to the external excitation process. Anisotropy and inhomogeneity are introduced through the local parameter values $\Gamma(x, y) \equiv K_{c\parallel}/K_{c\perp}$. Further details concerning the background and numerical implementation of the model (MATLAB encoded) may be found in [198]. In particular, as explained in that reference, the critical current parameter of Nb is better described as a function of

the local field, i.e. $K_c = K_{c0}[1 + B/B_0]^{-2}$. In this work, this property acquires relevance because flux is purposely concentrated in particular spots. In fact, such dependence produces more realistic predictions of the measured electromagnetic quantities.

The reconstruction of the experimental current-density streamlines from the measured $B_z(x, y)$ maps, used below to validate the simulations, follows the FFT-based Biot–Savart inversion described in Section 3.3.2 and Box 3.2. The superconducting film thickness and the sample-to-indicator distance, which set the inversion kernel, were 200 nm and 10 μm , respectively. As the superconducting sample covers most of the imaging area, the images were numerically extended by padding them with the value of the applied field so that the total image size is three times the original one. This leaves a region 3 mm wide where $B_z = \mu_0 H_a$ all around the sample, discouraging the inverted current lines from flowing through the periodic boundaries created by the Fourier-based technique and thereby ensuring that the current lines form closed loops.

7.4 Results

7.4.1 Tailoring Nb thin films: laser-processed meta-surfaces

Taking advantage of the observed anisotropy of the critical current density $\mathbf{J}_c$ for these laser-patterned samples [198], we designed and fabricated proof-of-concept structures aimed at controlling the magnetic flux penetration in the superconducting state (Fig. 7.2). In this figure, we sketch the line-wise meandering laser scanning protocol and the analyzed laser-patterned designs. The targeted orientation of the LIPSS domains is obtained by commanding the laser beam polarization direction (Fig. 7.1). As shown in Fig. 7.2a–c, basically parallel groove nanostructures are achieved which, on average, display a lateral spatial periodicity of $\approx 260\text{--}270\text{ nm}$, and a depth (from crest

to trough) of $\approx 80 \pm 15$ nm. Further details on the experimental characterization of the emerging LIPSS spatial frequencies, and on their theoretical simulation from the optical properties of Nb, are given in the supplementary materials of Ref. [77].

Conceptually [198], by using LIPSS, one might divert the natural tendency of the magnetic flux penetration, otherwise strongly influenced by demagnetizing effects in 2D geometry. The tested featured designs are (by order of complexity): (i) the *fringe structure* in Fig. 7.2d, aimed at the channeling of magnetic flux along the central line of pristine Nb, (ii) the *cross geometry* (Fig. 7.2e) aimed at “compensating” the preferential penetration of magnetic flux along the center of the film sides, and (iii) the *top-hat geometry* (Fig. 7.2f), a specific pathway for the penetration of magnetic flux, constructed by defining a structure of magnetic domains on the film surface. This design aims to “guide” the penetration of flux along a tailored path, which is flanked by parallel ridges of the LIPSS in the neighboring laser-patterned domains (expectedly confining the flux dynamics within the path by hindering the penetration perpendicular to the LIPSS). For brevity, Fig. 7.2 shows specific instances for each geometry, though equivalent designs with different sizes were also tested.

7.4.2 Routing magnetic flux penetration

Magnetic flux channeling and diversion

In Fig. 7.3, we present direct imaging of the magnetic field landscape obtained by the MOI technique, demonstrating the diversion of flux penetration induced by texturing the sample with domains of properly oriented LIPSS. Images correspond to two values of increasing out-of-plane applied magnetic fields after zero-field cooling (ZFC) at a stable temperature of $T = 8$ K for the fringe and cross-shape geometries. According to SQUID magnetization measurements, this temperature selection guarantees the restraint of the unstable flux avalanching processes [237] that were shown in Chapter 5 to be a prominent effect

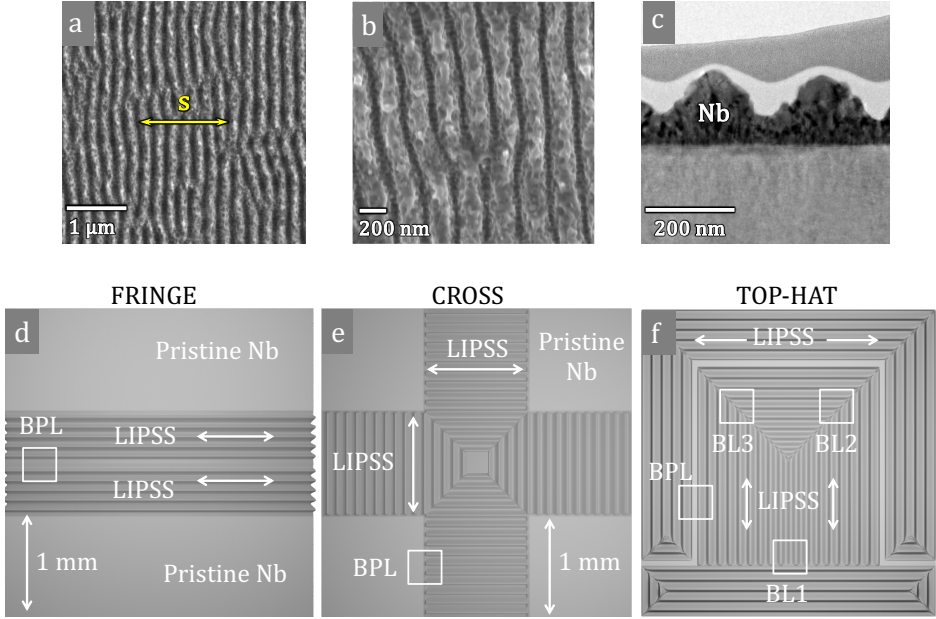


Figure 7.2: Panels (a–b) show top-view SEM images of a LIPSS domain observed by secondary electrons and in-lens detector, respectively, and (c) a cross-sectional TEM image. The yellow double arrow indicates the laser polarization. (d–f) show the geometries investigated in this work. The domain structures in (e) and (f) were created by programming the beam trajectories and rotating the sample by an angle of 90° , according to the pre-established configurations as explained in Fig. 7.1. In the three panels, plain gray corresponds to non-irradiated areas while striped areas are indicative of the laser-scanned parts. The transition between pristine and irradiated zones, as well as the boundaries separating the areas with different orientation of the LIPSS, are indicated. Small rectangles point out boundary regions that were investigated in detail (as shown in Section 7.5): BPL is for Boundary between Pristine & LIPSS regions and BL for Boundary between LIPSS domains.

well below T_c .

We notice that the *fringe* geometry displays a noticeably enhanced flux penetration along the central channel (Fig. 7.3a,b), even beyond the distance reached in the pristine region. In fact, as shown in Fig. 7.3b, the LIPSS act as a barrier for the penetration of magnetic flux. On the other hand, the transversal LIPSS structure in the *cross-shaped* geometry “protects” the central part of the sides from the usual enhanced flux penetration across that spot (Fig. 7.3c,d). These features have been repetitively observed in a wide range of temperatures (see the supplementary materials of [77]). In order to understand the physical origin of the experimental behavior observed in MOI, we carried out simulations of the penetration of magnetic flux. To that end, we put forward theoretically predicted current density streamlines and related magnetic flux profiles that capture the principal features revealed in the experiments (Fig. 7.3,i-l). In brief, numerical simulations have been performed by a CS formulation [230, 231], that takes into consideration a non-homogeneous pinning landscape (as described in Section 7.3). This coarse grained electromagnetic model uses a conduction law whose main parameter is the so-called local critical current density J_c (maximum non-dissipative current density allowed by the material at a given sample location). In terms of J_c , the laser treated areas are simulated by using an anisotropy factor Γ , as in Eq. (7.3), i.e. the ratio between the critical current density flowing parallel ($J_{c\parallel}$) and perpendicular ($J_{c\perp}$) to the LIPSS, and relative to the pristine Nb value. As sketched in the plot and justified below, here we have used $\Gamma = 3$, i.e. $J_{c\parallel} = 3J_{c\perp} = 3J_c$, with J_c the original value of the isotropic critical current density, characterizing the untreated Nb film regions.

We note that our results indicate that the critical current density is increased along the LIPSS (i.e. parallel to the LIPSS ridges), and keeps its original value perpendicular to them. We also point out that the ratio $\Gamma = 3$, that produces good agreement with the MOI observations, was obtained from the experimental discontinuity-line criterion applied to our MOI experiments in saturated conditions (see

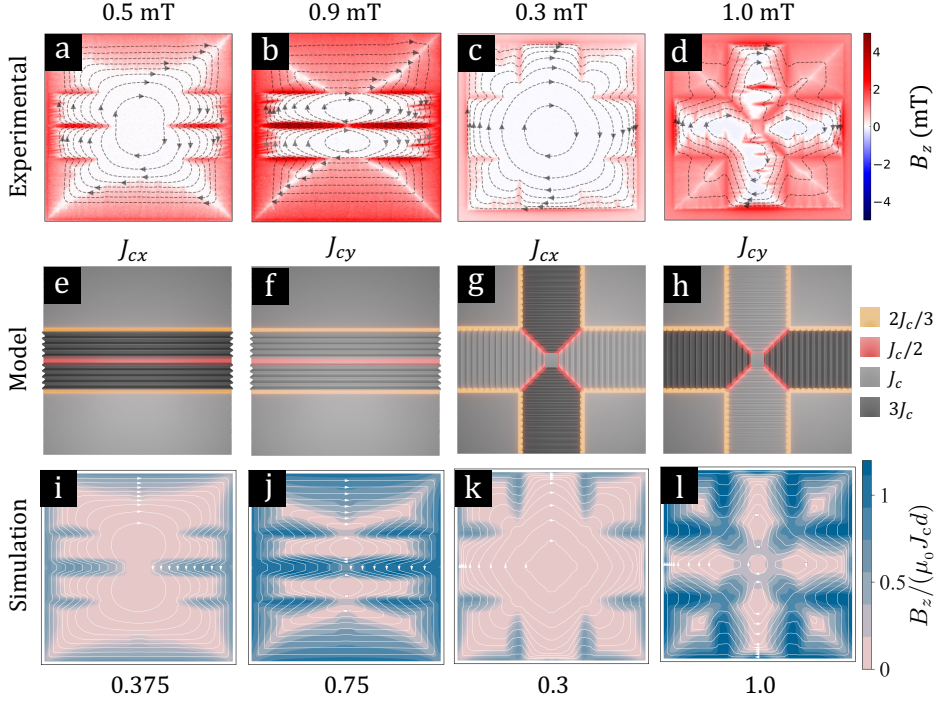


Figure 7.3: Magnetic flux penetration and current density streamlines in a laser patterned Nb film with the tailored *fringe* and *cross-shaped* structures of LIPSS parallel to the sample's perimeter. The experimental MOI images (a–d) were obtained at the temperature $T = 8\text{ K}$ and for the out-of-plane applied fields $\mu_0 H_a$ (mT) indicated on top. The reconstruction of the current density streamlines based on the inversion of the recorded magnetic flux density on top of the sample B_z is also shown. Arrowheads indicate the flow direction of the streamlines. The intermediate panels (e–h) sketch out the anisotropic and inhomogeneous current density model used in the simulations below. The lower panels (i–l) show theoretically predicted current density streamlines induced in the film obtained from the critical state model (see text), together with the related density plots of $B_z(x, y)$. Normalized dimensionless units are used both for the applied field and the resulting flux density (J_c is the critical current density and d is the thickness of the sample).

the supplementary materials of Ref. [77]).

Finally, it must be underlined that, in order to reproduce the experimental observations, a reduction of the critical current density within the boundary regions between patterned/unpatterned regions has to be assumed (see Fig. 7.3e–h). In simple terms, the enhanced flux penetration along the boundaries between the pristine and laser treated areas is concomitant with the sag in the current density streamlines. Such effect denotes a decrease in the modulus of the circulating current density J because, in steady state conditions, the separation between the streamlines is inversely proportional to J .

Notoriously, the phenomenological CS model (in particular, its simplified numerical template) captures the main features of the experiment, i.e.

- (i) Magnetic flux penetration (darker areas indicate higher flux density) is preferential along the boundaries in between patterned and unpatterned regions.
- (ii) In the *fringe* structure, flux penetrates faster along the central channel, followed by the penetration along the lateral boundaries and the non-treated zones, and, eventually, along the central part of the treated areas.
- (iii) In the *cross* structure, the flux density on top of the untreated square regions (corners of the sample) develops a deformed pyramidal roof structure, with edges deviated from the diagonal of the square.
- (iv) Eventually, at higher magnetic fields, the magnetic flux reaches the central area of the cross, by entering along the diagonal boundaries between the domains.
- (v) The flux density profiles within the patterned domains of the *cross* develop a fish-like silhouette (Fig. 7.3d,l) that reveals the underlying anisotropy, i.e. $B_z(x, y)$ penetrates faster parallel to the LIPSS.

Summing up this section, flux dynamics has been tailored by the underlying pattern of isotropic/anisotropic LIPSS-covered areas and their boundaries. Specifically, the natural trend in square platelets in transverse field, i.e. the inhomogeneous penetration of magnetic flux, mostly across the central part of each side [235], has been modified. In the laser treated samples, flux is driven along the interfaces of the areas patterned with LIPSS, that operate as principal flux distributors, eventually followed by channeling along the LIPSS themselves when flux enters the laser treated areas.

Steering magnetic flux: fabrication of pathways

Taking advantage of the above observations, we designed a surface topography that combines neighboring laser treated and untreated domains, acting as a magnetic hose for the flux penetrating in the film, i.e. the top-hat structure sketched in Fig. 7.2f. Recall that the targeted “magnetic path” combines the enhanced penetration along the boundary lines and the hindered drift of flux across the flanking ridges defined by the LIPSS.

We first concentrate on the laser-patterned surface domains. A close inspection by SEM of the boundary regions indicated in Fig. 7.2f is presented in Fig. 7.4. The microstructural detail around the channel (Fig. 7.4a, region BPL) reveals a $\approx 30\text{ }\mu\text{m}$ -wide band of, allegedly, original Nb surrounded by a transition zone, that eventually connects to the nanostructured area (the targeted gap between the centers of the laser beam at the channel sides was $50\text{ }\mu\text{m}$). The transition band ($10\text{ }\mu\text{m}$ -wide) displays an induced nanoscale roughness but no pronounced LIPSS, which in fact start progressively forming beyond this point and are aligned parallel to the channel. Here, it is important to recall that in this case the laser beam scanned the sample perpendicular to the channel, in a meandering bidirectional trajectory. Therefore, the transition boundary is formed by the superposition between the successive beginning/end points of the scan lines, and may be affected by power delivery and beam acceleration issues related

to the smaller overlap of the laser spot energy density profiles.

The morphological characteristics of the transition are similar to those of the BPLs in the cross sample, but wider in that case (see the supplementary materials of [77]). The difference stems from the fact that, owing to the targeted arrangement of the LIPSS, for the latter case the laser scan was parallel to the boundaries and the overlap was smaller owing to the elliptic shape of the spot (see Section 7.2.1). However, in both cases, the widths of the transition zones are of the same order and always smaller than the semi-axes of the elliptic Gaussian laser beam.

In fact, as described in Ref. [198], the formation of well-structured LIPSS is a two-fold process strongly dependent on the initial profile of energy density deposited by the laser beam and on its superposition with the successive pulses in the scanning process. Consequently, a well-defined boundary requires a precise overlapping strategy compromising the finite beam overlap and the precise allocation of the sample's rotation axis for the case of complex LIPSS-patterned domain geometries. With our setup, a mismatch of about 10–15 μm was achieved, and better and better results were obtained by a “trial and error” strategy positioning the sample on the holder, limited by the positioning and repeat accuracy of the motorized stages. More advanced laser-scanner technology with full (“single-spot” (pixel)) control of polarization direction, scanning direction and scan velocity is envisaged for the future.

The transitions of the surface topography in between domains with different LIPSS orientations (spots BL1, BL2 and BL3) are detailed in Fig. 7.4b–d. In the regions identified as BL1 and BL3, LIPSS of the neighboring domains intersect around the orientational boundary dividing them, as suggested by the highlighted discontinuous lines in the images. On the contrary, in the zone BL2 there is a gap between them, with a transition band similar to the one shown in Fig. 7.4a. By comparing SEM and MOI images of this sample, the influence of the different morphological characteristics of the intersections between LIPSS domains on their superconducting behavior will be analyzed.

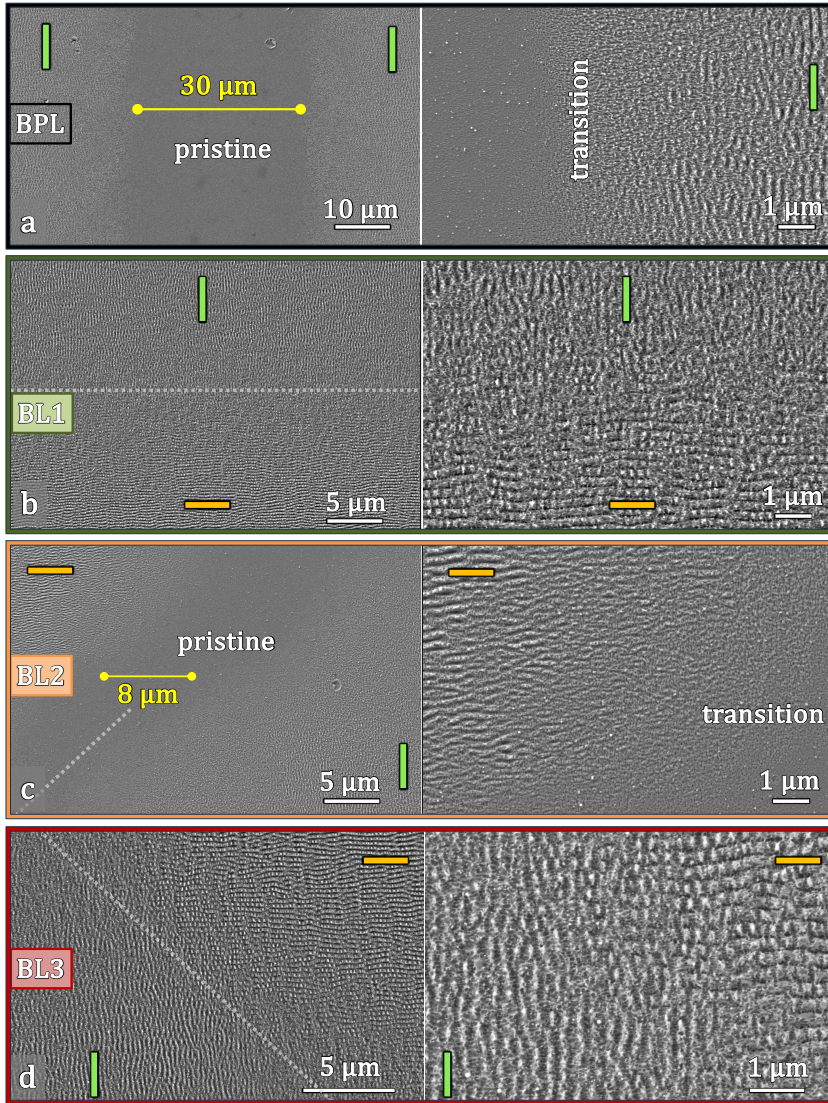


Figure 7.4: SEM images of the boundary area BPL (a), BL1 (b), BL2 (c), BL3 (d) marked in Fig. 7.2 for the top-hat geometry. For each case, different areas with different magnifications are displayed. Dotted white lines are a guide to the eye for better visualization of the laser-patterned domain boundaries. The colored bars indicate the microscopic orientation of the LIPSS (orange: horizontal, green: vertical).

Fig. 7.5 shows a selected set of MOI images obtained for the top-hat path upon increasing applied magnetic field from 0 to 1.2 mT after ZFC down to 8.28 K.

Initially, flux enters the sample following the designed path, symmetrically from both sides. In fact, by closer inspection of the MOI images around the path (see inset in Fig. 7.5 at 0.3 mT), it is noteworthy that the field is higher at the very boundaries than at the central part. Then, at applied fields of 0.3–0.4 mT the magnetic flux pattern deforms properly, following the drawn path. However, against expectations, upon further increasing the field, some flux leaks in at the separation between some neighboring domains and eventually leaks along the LIPSS. Leakage takes place:

- (i) partly along the initial penetration direction, thus following the boundary line (BL1) that separates two LIPSS domains (aligned vertically/horizontally in its upper/lower side).
- (ii) for fields higher than 0.6 mT, flux also enters along the upper-right large diagonal (BL2), again defined by boundary-lines between orientational domains (Fig. 7.2).

Notwithstanding the above, the symmetrically disposed upper-left diagonal (BL3) shows a perfect behavior; that is, the induced supercurrents in this area of the sample allow appropriate shielding of the magnetic field, avoiding leaks from the designed flux channel.

From the above, one may conclude that well-defined boundaries are necessary in order to achieve optimum channeling conditions (as revealed by comparison of the BL2 and BL3 zones), together with some further considerations as deduced from the non-perfect behavior in BL1. In Section 7.5, the flux penetration properties observed by MOI and described phenomenologically by the CS theory will be further connected to the laser-induced microstructural properties through complementary experimental information.

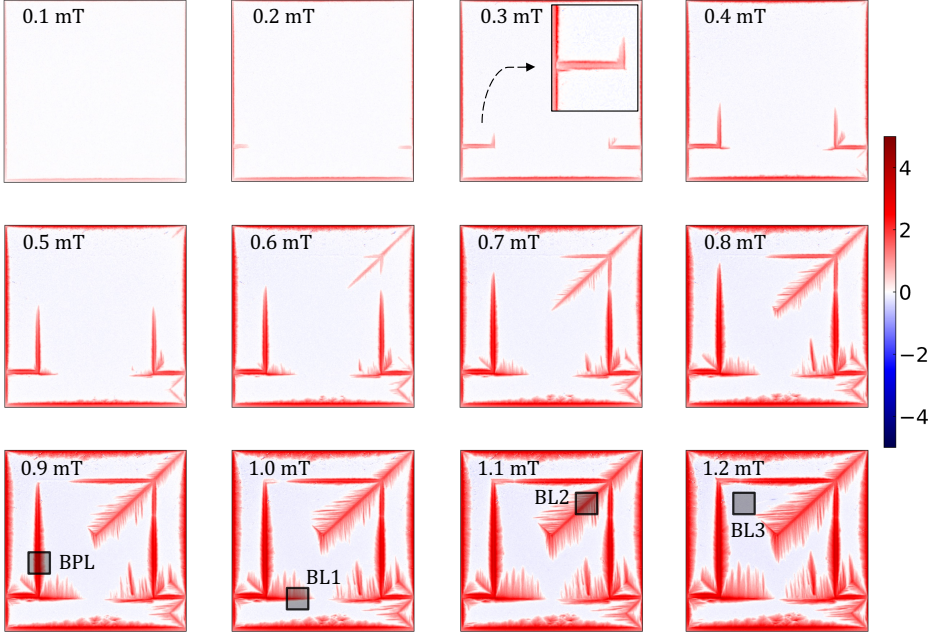


Figure 7.5: Magneto-optical images of the flux penetration process in the top-hat channel described in Fig. 7.2f. The increasing values of the applied magnetic field are indicated on the successive steps. The thin film was zero-field-cooled to the base temperature of 8.28 K and field subsequently ramped up. Darker areas correspond to the higher values of the local magnetic field (basically along the top-hat channel), and are accompanied by ramifications along the neighboring LIPSS in the boundaries between orientation domains (compare to Fig. 7.2f). BPL, BL1–BL3 outline areas around the boundaries between domains, with the presence of LIPSS with different orientations that have been investigated in detail.

Parceling magnetic flux

To this point, we have described the behavior of laser patterned metastructures under monotonic increasing field processes for the geometries sketched in Fig. 7.2. Below, we reveal a new feature emerging under alternating field cycles. More specifically, we show that by applying low frequency, i.e. $\lesssim 100$ mHz AC cycles with decreasing field amplitude, one may fabricate and store flux packets of controllable size, consisting of alternating units of positive and negative flux (see Fig. 7.6).

As shown in the upper panels, by virtue of a few number of cycles, we achieve a rather robust structure of alternate positive/negative flux segments within the channel, of size Δ controllable by the rate of decrease of the applied field amplitude δH_a . We notice that the formation of these chains is rather independent of the actual environment surrounding the flanking LIPSS domains.

Again, from the physical point of view, the construction of these regular alternating flux density profiles also finds an explanation in the CS theory [230, 231]. As shown in Fig. 7.6g,h, even the simplest formulation (Bean’s model for infinite slabs in parallel magnetic field) predicts a “saw-tooth” profile of positive and negative flux bits when the applied field is cycled along the process $H_1 \rightarrow H_2 \rightarrow H_3 \rightarrow H_4 \rightarrow H_5$. For the more realistic 2D geometry accomplished in our experiments with thin films, the structure within the channels flanked by LIPSS domains is provided with high fidelity by the numerical solution of the CS model as displayed in the panel i of Fig. 7.6.

7.5 Discussion

The results presented so far provide evidence of the successful achievement of magnetic flux diversion and channeling on superconducting Nb thin films by an appropriate laser surface nanostructuring, in particular as concerns the role of orientation domain boundaries. As shown in Section 7.4.2, the actual microstructural modifications in-

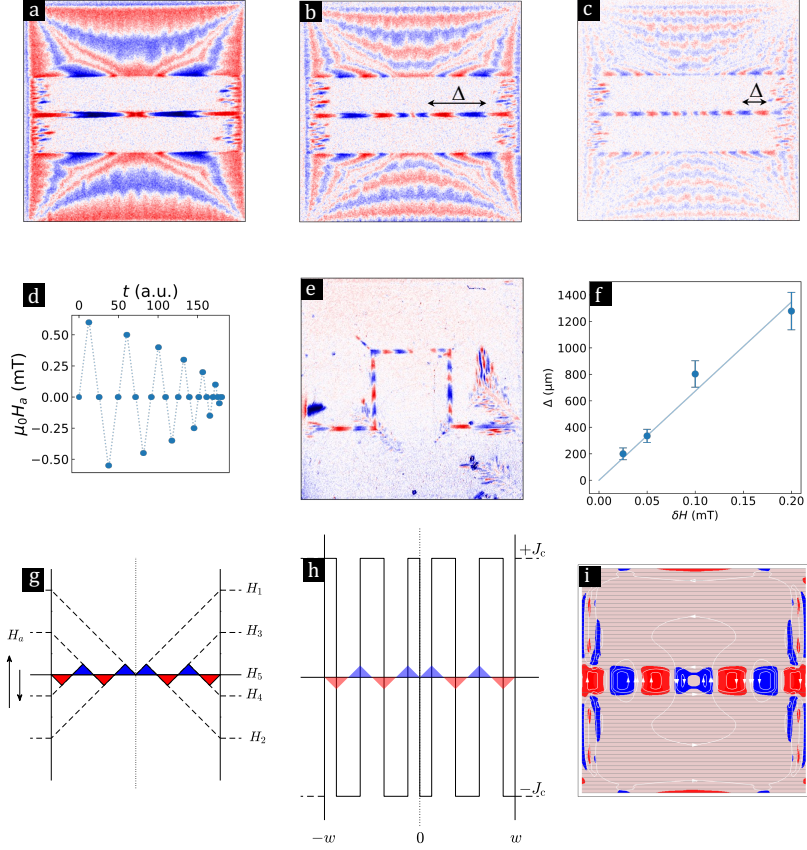


Figure 7.6: Magneto-optical images of the *flux parcelling* process along the channel of the *fringe* geometry (a–c) for different snapshots corresponding to the “magnetization” protocol shown in panel (d). (e): the same for the *top-hat* geometry. In the sequence (a–c), we used $T/T_c = 0.776$, and reductions of amplitude $\mu_0 \delta H_a = 0.2, 0.1, 0.05$ mT respectively, whereas panel (e) corresponds to $T/T_c = 0.91$, and $\mu_0 \delta H_a = 0.05$ mT. The same color bar as in Fig. 7.5 was used. In panel (f) we plot the experimental “flux packet” size Δ as a function of $\mu_0 \delta H_a$. Panels (g)–(i) depict the supporting theoretical framework. (g),(h) display the flux penetration and current density profiles induced by a demagnetizing process under Bean’s model 1D approximation, while (i) shows the expected flux density and current density streamlines obtained through a numerical 2D model for a straight channel structure.

duced by the laser at the inter-domain frontiers are crucial for the eventual behavior of magnetic flux. In particular, flux enters preferentially along the rough transition regions between pristine Nb and the domains of well-ordered LIPSS (BPL in Fig. 7.2f), and may be conveniently expelled from those areas with neat transitions between different orientations of the LIPSS (BL3 in Fig. 7.2f). This second condition is nevertheless not sufficient, as one may notice by the flux leakage along the boundary BL1 in Fig. 7.5. As a rule of thumb, the “clean transition” requirement must be adjoined with a second basic condition: within the flux-free areas, shielding current loops should be allowed to flow with ease along shape-preserving paths following the configuration of the channel. This basic property of every flux screening distribution may be assessed from the observation of the current density loops in Fig. 7.3, together with the improvement of the behavior in the region BL1 by defining an oblique boundary, as it is noticeable comparing Fig. 7.5 and Fig. 7.6e.

In order to further investigate the properties of the above mentioned “rough” transition zones that restrain the flux penetration channels, additional experiments collected in Fig. 7.7 were performed. Fig. 7.7a,b shows the MOI intensity decay with temperature measured in different zones of the fringe sample after field cooling to the base temperature under an applied field of $\mu_0 H_a = 4 \text{ mT}$ and subsequently turning the field off. This allows an estimation of the local critical temperature, defined as the temperature where MOI relative intensity reaches the constant value beyond the transition. As compared to the value of the non-irradiated zones in the same sample, $T_{c,\text{pristine}}$, we observe a small decrease in the LIPSS-patterned regions, $T_{c,\text{LIPSS}}/T_{c,\text{pristine}} \approx 0.975$, which further reduces to $T_{c,\text{transition}}/T_{c,\text{pristine}} \approx 0.95$ in all BPL boundaries. On the other hand, cAFM measurements at room temperature across the magnetic flux channel of the same sample (Fig. 7.7c) indicate significant changes on the measured contact resistance by tapping against the region around the path.

The significant variations of the local electrical conductivity enable

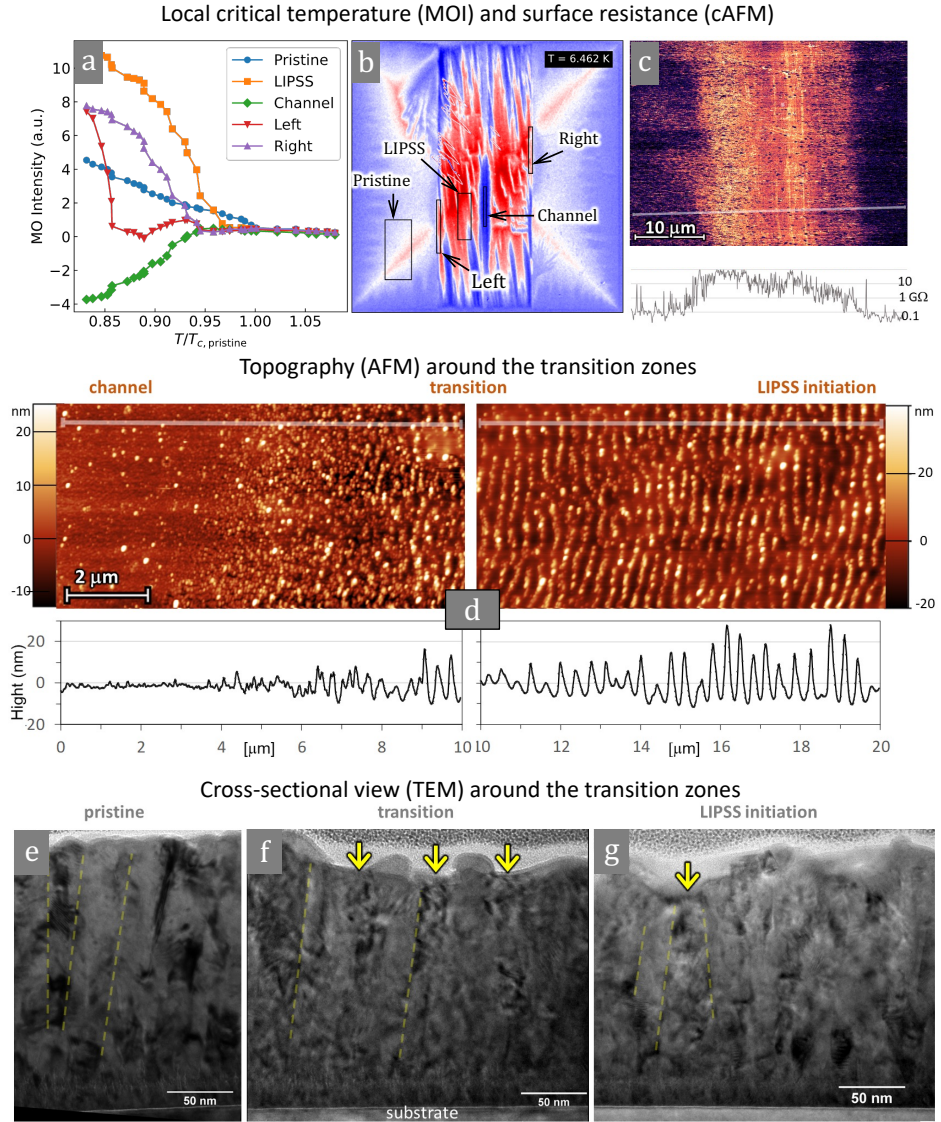


Figure 7.7: (a) MOI intensity decay as a function of temperature at different zones (highlighted by squares in panel (b)), (c) cAFM image at room temperature around the magnetic flux channel of the fringe sample. The graph shows the resistance measured along the line drawn in the image above, (d) AFM images and corresponding line profiles near the flux channel of the top-hat sample (consecutive scans over $10 \times 3.5 \mu\text{m}^2$ areas). TEM images (cross-sections) of a non-patterned area (e) and BPL zone (f)–(g) of the cross sample. Dotted lines are a guide to the eye. Yellow arrows point to some grains featuring larger number of dislocations and defects.

us to use the cAFM technique for mapping the different zones of the analyzed samples with excellent correspondence to the magnetic flux channels observed at low temperatures.

The topographical analysis around the *magnetic flux channels* (see Fig. 7.7d) reflects a progressive increase in surface roughness from the channel in agreement with SEM images. It must be mentioned that patterning occurs gradually, with an undefined spatial period in the transition area until, eventually, self-organization initiates and gives rise to the LIPSS. Note that the ripples near the boundary have significantly lower depths than those inside the LIPSS-patterned regions. Remarkably, the topographical changes observed by AFM around the magnetic channel are an indication of underlying changes related to the observed reduction of T_c and implied decrease of J_c observed by MOI. Whether these surface properties are the only factor affecting superconductivity at the local level deserves further investigation. For this purpose, the observation of cross-sectional views of the full Nb layer at the regions of interest was addressed. Reference lamellas were prepared from the transition zone in the *cross-shaped* sample for TEM observation.

TEM images of the different zones are shown in Fig. 7.7e–g. In all cases, we observed a perfect pairing between the Nb film and the substrate, as well as the absence of cracks, thus ruling out detachments or mismatch as the origin of the depleted superconducting properties. We notice that the pristine area shows that the Nb layer is composed by very small grains (size about 20–30 nm) which are basically disposed in a columnar arrangement. In the transition zone (Fig. 7.7f), there are small protrusions at the surface (in consonance with SEM and AFM observations) and the film thickness decreases slightly (by 5–10%). The TEM images of this area show that the influence of laser is not restricted to the surface or to a very thin layer below it. On the contrary, they reflect an increase in the number of dislocations and defects (larger number of darker lines in the images), specially in those grains pointed out by yellow arrows in the figure, in between the small surface protrusions. The columnar disposition of grains is still

clearly visible. In the transition region close to the LIPSS (Fig. 7.7g) the disorder and number of defects further increases and the grain columnar arrangement becomes less clear. Note also that there is a small zone of ≈ 25 nm in depth, close to the Si-substrate, where the grains are smaller. This zone remains apparently unaltered in all the analyzed regions, even when the LIPSS are fully formed. It is suggested that the important increase of defects in the grains of this transitional band, which also features a nanoscale rough topography on the surface, has an important influence on the superconducting properties, resulting in a decrease of T_c and J_c values (in the low magnetic field range used in MOI experiments).

7.6 Conclusions

The possibility of arbitrarily steering the penetration of magnetic flux in superconducting chips through the control of surface morphology has been demonstrated. Taking advantage of ultrashort laser processing pulse technologies, we have tailored surface domains consisting of oriented groove structures with submicron periodicity on Nb type-II superconducting thin films. By irradiating the samples with femtosecond laser pulses, we have found an expeditious, scalable and ecologically friendly physical procedure that enables the formation of metasurfaces useful for diverting, guiding or parceling magnetic flux with density in the range of millitesla. Our main findings are:

- The orientation of surface ripples (LIPSS, type LSFL) within a given domain produces an anisotropic magnetic field penetration. In our case, the critical current density along the ripples is enhanced by a factor $\Gamma = 3$, as compared to the perpendicular flow. This factor is obtained experimentally by comparing the field penetration parallel/perpendicular to the arms of the cross structure in Fig. 7.3, and has been used as a simulation parameter that reproduces the experimental flux penetration landscapes in a wide range of experimental conditions. As the

physical counterpart, magnetic flux becomes trapped in the valleys. Pictorially speaking, flux tubes as well as supercurrents are guided along the grooves of the ripples, acting as rails, and are restrained for perpendicular displacements.

- The boundaries between neighboring LIPSS orientation domains as well as those between them and the adjoining pristine areas play a major role in the penetration of magnetic flux. As shown in Figs. 7.5 and 7.6, precise switching of the magnetic flux penetration is achieved by defining neat flanks around the flux path as well as between orientation domains in the targeted flux protected areas. However, fuzzy transition zones in the internal boundaries between the domains act as magnetic flux scatterers, since the flux easily glides along these zones, eventually spreading in the direction of the closest ripple structures.
- Alternating field operation has been tested and increases robustness against imperfections in the focused topography. It is also remarkable that stable and tunable flux packets within the flux pathway are obtained by taking advantage of the flux pinning property, together with the flanking barrier effects.

In order to achieve the targeted surface morphology, some critical processing parameters in the laser treatment must be considered:

- The irradiation conditions must be adjusted in a typically large spot overlap of adjacent laser irradiations so as to ensure the material reconfiguration induced by the selective energy deposition [198]. This implies a trial and error material-dependent process ending up with a set of parameters (laser power, pulse repetition frequency, scanning speed and overlapping between different scan lines). Optimized values are given in Section 7.2.1.
- Magnetic flux hosing effects are improved by a combination of enhanced superconductivity in the LIPSS domains and reduced

parameters in the boundaries with untreated zones. This (initially unwanted) effect relates to the material nanostructure induced at the end points of the laser beam scanning lines. In these zones, we have observed an increased number of defects, disorder, a noticeably higher resistivity at room temperature, and clearly decreased superconducting critical temperatures T_c as deduced from TEM, AFM and MOI local probes, respectively. This picture is also consistent with a reduced critical current density J_c for consistency with numerical simulations of the overall flux penetration features.

- The neat transition across the boundaries between orientational LIPSS domains is crucial to allow supercurrent percolation and flux expulsion within the protected areas. For this purpose, a subtle combination of laser scanning conditions is necessary, compromising the local energy profile of the beam and the superposition of successive laser pulses over a given area. Well-defined laser-patterned domain boundaries require good overlapping among laser scans in different directions. At the same time, care must be taken so as to avoid overtreatment, that ends up with holes in the Nb film, and as a consequence with unwanted flux penetration.

Finally, we recall that this fast (mm^2/min) and scalable laser-assisted technology for tailoring thin-film surfaces with selective magnetic field penetration may find applications in several fields such as the enhanced shielding of superconducting chips by framing the protected area with a band of LIPSS parallel to the surface, the functionalization of monolithic superconducting metamaterials, vortex diodes, guides, etc., or even the accomplishment of tunable superconducting/magnetic hybrid structures [238], which may find application in superconducting microelectronics. In the specific context of this thesis, the LIPSS metasurfaces demonstrated here offer a complementary route to the magnetic-field resilience strategies discussed for the CPW resonators of Chapter 5: a peripheral band of properly oriented LIPSS

could, in principle, frame the sensitive resonator area and divert the penetrating flux, together with the accompanying thermomagnetic avalanches, away from the high-RF-current edges.

Chapter 8

Conclusions and outlooks

This thesis sets out from a single premise: map locally the physical parameter governing the phenomenon under investigation, and subsequently confront the obtained map with the averaged observable returned by a conventional global measurement. The parameter to be mapped changed from chapter to chapter, from the magnetic flux density to the screening currents that determine its penetration and then to the radio-frequency current of the resonant mode. The approach remains the same. Each experimental chapter was an instance of that methodology, and each answered one of the questions posed in the introduction.

8.1 Summary of the main results

The instruments. Part of the work performed during this thesis consists of instrumental developments rather than measurements. Both, technical accomplishments and data collection are reported together since it is not always straightforward to separate them from the results that follow: each of them constituted a milestone on the quest to remove a specific obstacle that made one of the questions of the introduction unanswerable, and none was undertaken for its own sake.

The first question could not even be posed properly. Deciding whether the RF drive provokes avalanches requires monitoring the magnetic flux and the resonance simultaneously, since otherwise avalanche activity has to be inferred from the very microwave trace whose interpretation is in question. The magneto-optical setup was already in place, but its indicator carries a metallic mirror in which induced eddy currents both load the resonator and brake the avalanches themselves, suppressing precisely the delayed events that had to be counted. Removing that mirror, and adapting the optical cryostat to accommodate a microwave line, turned an open question into an answerable one.

The second question was hindered by the absence of the specific tool able to address it. Identifying which regions of a resonator govern its resonance frequency requires mapping the radio-frequency current of the resonant mode, and no instrument capable of doing so existed in Liège when this work began. The scanning laser microscope presented in Chapter 4 was therefore designed, assembled, programmed and calibrated over the course of this thesis, with that precise objective in mind. Resolving the sample's response in time raised a second obstacle, since a static map hides the dynamics entirely. In order to lift this obstacle, the multiharmonic scheme was developed on that platform.

Revealing the erratic RF response produced by avalanches.

Chapter 5 combined magneto-optical imaging with microwave transmission on NbTiN coplanar-waveguide resonators, first sequentially and then simultaneously. The sequential measurements showed that the erratic response of the resonance frequency to an applied field is not noise but the superposition of four distinct regimes: magnetic field screening, a magnetic perforation of the ground plane, avalanche-dominated penetration, and a smooth flux removal terminating in anti-avalanches. The threshold field for avalanche activity was extracted consistently from the magneto-optical images and from the microwave trace alone, which establishes the converse of the thesis's

usual direction: a superconducting resonator is itself a sensitive, non-invasive probe of thermomagnetic instabilities in its own film. This is an interesting finding which could permit, for instance, revealing the presence of flux avalanches in Al films (or other low- T_c superconductors used in technologically relevant resonators) for which no sub-Kelvin MOI has been developed.

Adapting the experimental setup to acquire both signals at once, in the mirrorless configuration described above, sharpened this correspondence into an event-by-event one and yielded three results that sequential measurements could not reach. First, it settled the long-standing question of whether the RF drive stimulates avalanches. Within the linear Campbell regime, where the drive shakes vortices without reconfiguring them, the effect is real but marginal. The drive does not create avalanches, it only precipitates them: a stronger RF amplitude breaks the metastability earlier, which is why higher powers produce more frequent but smaller jumps. No excess activity appears where the RF current is largest, suggesting that an otherwise stable configuration is not forced to enter an avalanche regime by the drive alone. Second, tagging individual avalanches to individual frequency jumps exposed the non-local character of the electrodynamics, since an avalanche in one resonator, or in the feedline, shifts the resonance of every resonator on the chip. This has a direct design consequence: reinforcing pinning only where the vortex density is largest is at best a partial remedy. Third, electromagnetic simulations showed that the frequency shift produced by an avalanche scales with the square of the local sheet current density and is governed by the avalanche's contact length with the conductor edge rather than by its area, which identifies the edges as the regions to protect.

Identifying which regions of a resonator matter, and when. Because avalanches are difficult to trigger in high- T_c materials, YBCO constitutes a natural alternative platform to conventional superconductors. Chapter 6 accordingly turned to YBCO resonators, while mapping a different local parameter: no longer the flux or the currents

that screen it, but the radio-frequency current of the resonant mode itself. In oxygen-deficient films, persistent photodoping was shown to be slow, stable and non-diffusive, in sharp contrast with the fast and highly diffusive bolometric response that is initially observed due to laser heating. Exploiting that contrast, the microscope built for this purpose mapped the local RF current density and resolved the standing waves at both the fundamental resonance and the second harmonic, identifying the regions that most strongly govern the resonance frequency. Because photodoping is reversible and spatially selective, it also constitutes a way of *modifying* a resonator optically, in contrast to the enhancement obtained by heavy-ion irradiation which is permanent and precludes any subsequent comparison on the same device.

The multiharmonic scheme then extended this mapping into the time domain: collecting several hundred harmonics of a chopped photoresponse with parallel lock-in amplifiers and inverting the Fourier series pixel by pixel yields a dynamic response with few-microsecond resolution, only limited by the highest measurable harmonic. The time-resolved maps showed that the notion of a sensitive region is more subtle than a static map suggests: the central conductor responds almost immediately, while weakly current-carrying regions of the ground plane contribute with a delay that grows linearly with distance and persists after the laser is switched off. The response is therefore neither instantaneous nor purely local, and the propagation of the modulation yields an effective thermal diffusivity of the substrate as a by-product.

Escaping the vortex losses dilemma by steering them. The conventional remedy against field-induced losses is to immobilize vortices with artificial pinning sites. However, this remedy sits in tension with the mechanism identified in Chapter 5, since pinning sites are known to favor avalanche triggering at low temperature. In Chapter 7, we showed that flux can be handled differently, thus taking the first step from diagnosis to design. Femtosecond laser pulses

imprint on niobium films a metasurface of oriented submicron ripples. Inverting the magneto-optical images into maps of the screening currents shows that these ripples raise the critical current density along their ridges by a factor of three relative to the perpendicular direction, rendering the film electrically anisotropic and the flux penetration correspondingly directional. The decisive finding, however, is that the ripples are not the principal actors: the boundaries between orientation domains are, acting as flux distributors along which penetration preferentially glides. Exploiting this hierarchy, flux was channeled along designed pathways, diverted away from protected areas, and parceled into stable packets of controllable size. Exclusion is thus not the only available strategy, and flux that cannot be kept out can be rerouted.

The common thread. Taken in sequence, these three chapters instantiate the trajectory announced in the introduction. Mapping the flux *revealed* the mechanism hidden inside an averaged signal; mapping the RF current *localized* the regions that are most vulnerable and exposed their dynamics; mapping the screening currents, and then patterning the surface, opened a route to *steer* the flux goes. They also converge on a single device concept: a superconducting resonator protected by the very laser-processing tools studied here and diagnosed by the very mapping techniques developed here. The remainder of this chapter reports a first, exploratory step toward that goal, then outlines the work that remains.

8.2 Preliminary results on a converging experiment: laser-machined resonators

The three techniques of this thesis meet in a single device concept: a coplanar-waveguide resonator carved directly into a plain niobium film by laser ablation, with no lithographic step at all. The long-term goal is a resonator whose surface is covered everywhere with LIPSS,

so that the flux-steering demonstrated on plain films in Chapter 7 is exploited across the whole device to prevent flux penetration and stabilize the resonance frequency over a wider magnetic field range. Laser machining is the natural route to such a device, but it comes with an unavoidable consequence: ablating the resonator geometry *inevitably* generates a narrow ($\sim 1\ \mu\text{m}$) band of laser-induced periodic surface structures along every etched border, and this band cannot be circumvented. Since the screening currents crowd precisely at those borders, the orientation of this border LIPSS relative to the meander, parallel or perpendicular, is expected to influence the flux penetration and the field resilience of the resonator from the outset.

This makes the logical first checkpoint not merely “can a resonator be machined?” but rather “does the orientation of the unavoidable border LIPSS already matter?” The study reported here is at exactly this stage: four devices sharing the same meander geometry, two with border LIPSS parallel (A-PAR and B-PAR) and two perpendicular (A-PER and B-PER) to the meander. The machining conditions were similar for all samples. Despite that, sample B-PAR did not present any clearly trackable resonance, and the results are consequently inconclusive at this early stage: the resonators cannot yet be machined reproducibly with consistent resonance frequency and quality factor.

On the encouraging side, some laser-machined devices do resonate, their temperature dependence is mutually consistent, and the DC resistance between feedline and ground plane exceeds $1\ \text{M}\Omega$, confirming that the ablation defines the resonator geometry cleanly. The best of the four samples (B-PER) exhibits a quality factor of around 200, a resonance frequency that can be tracked continuously under field, and the field hysteresis expected of a type-II superconductor, strongly hysteretic at low temperature and vanishing near T_c as pinning weakens. The temperature dependence of both the parallel and the perpendicular samples further indicates that a high critical temperature is maintained despite the aggressive machining. Laser machining, border LIPSS and all, can therefore produce a superconducting resonator of usable quality.

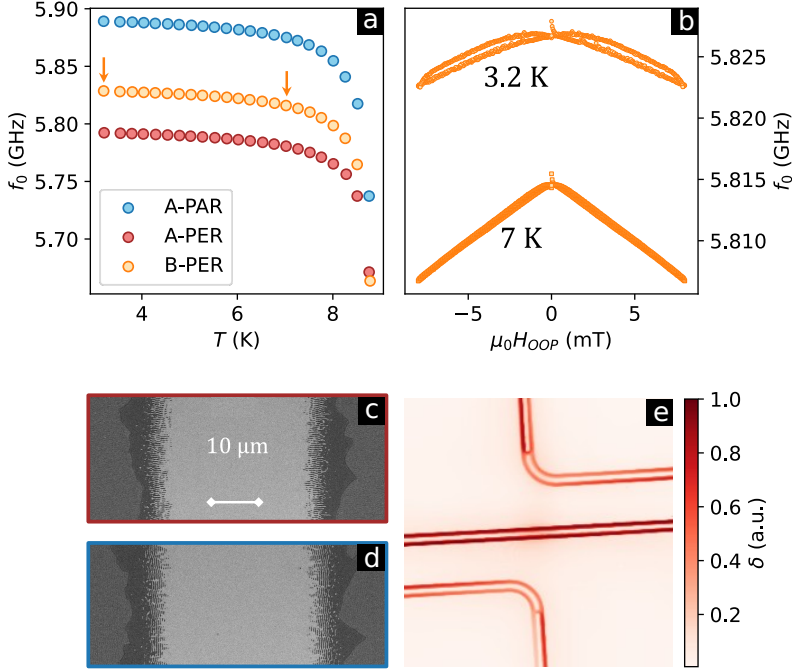


Figure 8.1: Laser-machined niobium coplanar-waveguide resonators. The three samples shown here are those exhibiting a trackable resonance, B-PAR being excluded for not meeting that criterion. All three share the same meander geometry, defined entirely by laser ablation of a plain Nb film, and differ only in the orientation of the unavoidable border LIPSS relative to the meander: parallel (PAR) or perpendicular (PER). (a) Temperature dependence of the resonance frequency, confirming that the devices resonate and behave consistently. (b) Resonance response under an applied magnetic field for sample B-PER. The hysteresis at low temperature and the quasi-linear dependence at high temperature are consistent with expectations. (c) (resp. (d)) SEM images of the central track of A-PER (resp. A-PAR), with LIPSS visible perpendicular (resp. parallel) to the border. The frame colors match the legend of panel (a). (e) SLM image acquired on sample A-PER. Here the gaps are brighter than the Nb because they absorb more light and hence heat up more.

The variability across the four nominally identical devices is, however, not yet understood, and it clouds precisely the orientation question the experiment was meant to answer. With functionality varying so much from device to device, the data do not allow the effect of the border-LIPSS orientation to be separated from ordinary fabrication scatter. Bringing this device-to-device spread under control is the prerequisite for answering the orientation question, and in turn for the ultimate goal of a fully LIPSS-covered resonator.

Scanning laser microscopy of the best-performing sample confirms that the mapping approach of Chapter 6 transfers to these low- T_c devices. The reflectivity contrast is reversed with respect to the YBCO case, because niobium reflects more strongly than the Si/SiO₂ substrate, so the conductor tracks now appear bright against a dark background, the substrate heating up more than the superconducting film. Despite this inversion, the carved meander is resolved sharply, demonstrating the versatility of the scanning laser technique for mapping RF currents in devices (see Fig. 8.1(d)).

8.3 Experimental limitations

The spatial resolution gained in this work is not obtained for free, and three limitations should be kept in mind when reading the results above.

The first is invasiveness. Magneto-optical imaging requires an indicator to be laid directly on the device, and Chapter 5 quantified the price: the indicator increases the effective permittivity and the dielectric losses, shifting the resonance frequency and lowering the quality factor. The simultaneous measurements are therefore performed on a device that is measurably perturbed by its own diagnosis, which may well mask a small RF-induced reduction of the avalanche threshold field. Removing the metallic mirror, as was done here, suppresses the eddy currents but not this dielectric loading, and the constraint is not specific to magneto-optics: any technique placing a solid sensing layer

on the device, quantum diamond microscopy [157] included, loads the resonator in the same way. Genuinely non-perturbative imaging of a resonator under microwave excitation therefore remains an open experimental problem.

The second is resolution. The micrometer resolution of both techniques is set by the optical wavelength and the sensor-to-sample distance, and it is coarse compared with the superconducting length scales, $\xi \sim 4 \text{ nm}$ and $\lambda \sim 400 \text{ nm}$ in NbTiN. Individual vortices are never resolved here and what is measured is always a coarse-grained flux density.

The third is interpretational, and it concerns the notion of a *sensitive region* itself. What SLM measures is a global frequency shift produced by heating a spot much smaller than the device, whose thermal profile nonetheless extends well beyond the spot. A weakly heated region carrying a large RF current and a strongly heated region carrying a modest one produce the same shift, and disentangling them is not possible from the map alone. The sensitivity maps of Chapter 6 are therefore best read as convolutions of the RF current density with a thermal kernel, and the effective parameters extracted from them, notably the thermal diffusivity, as effective rather than metrological quantities. This limitation also indicates where the interpretation can be improved. Modeling the heat diffusion in the film and the substrate, and coupling it to an electromagnetic simulation of the resonant mode, would give access to the thermal kernel itself. Deconvolving it from the measured maps would separate the RF current distribution from the thermal response, and would turn the diffusivity and the thermal boundary conductance into quantitative estimates rather than effective values. Developing such simulation tools is a natural continuation of this work.

8.4 Outlook

The microscope beyond this thesis. The scanning laser microscope has already been used to investigate systems well outside the scope of the preceding chapters, leading to two manuscripts in preparation.

The first application uses the microscope as a purely thermometric probe on YBCO multiterminal devices. At fixed bias current, the modulated local heating gives access to dR/dT pixel by pixel. Repeating the measurement across the transition and integrating numerically permits reconstructing the local resistive transition, and thus yields a map of the critical temperature itself rather than that of the current.

The second concerns micronetworks of silver nanowires, candidates for transparent conductive materials. The current does not spread uniformly over such a network; it concentrates along a subset of percolating paths. The electrical photoresponse identifies the wires that actually carry it, while reflectivity imaging reveals where the network physically fails. Following both during electromigration shows that the wires breaking first are those for which the current map had singled out beforehand.

Extending the mapping techniques. The multiharmonic time-resolved scheme of Chapter 4 is likewise not specific to superconducting resonators, and its time resolution is set by the highest measurable harmonic rather than by any mechanical scanning speed. Adding lock-in amplifiers in parallel, or improving the signal-to-noise ratio at high harmonic order, translates directly into finer time resolution, and the method transfers to any system whose global response can be modulated by a scanned local perturbation. On the photodoping side, reducing the laser spot size and increasing the scanning speed would turn the proof of concept of Chapter 6 into a practical technique for erasable laser-based scribing of superconducting circuits. The same persistent photoresponse exists in other cuprates, such as oxygen-deficient $\text{GdBa}_2\text{Cu}_3\text{O}_x$ [180] and probably $\text{Bi}_2\text{Sr}_2\text{Ca}_2\text{Cu}_3\text{O}_{10+x}$ [181], and analogous persistent photoconductivity in oxides such as stron-

tium titanate [182] and lanthanum aluminate [183] offers a natural extension.

From steering to protecting. On the LIPSS side, the open question raised by Chapter 7 is the microscopic origin of the boundary behavior: whether the reduced T_c and J_c observed at domain boundaries stem from the surface topography alone or from the underlying defect structure revealed by TEM. Answering it would turn the empirical recipe for flux guiding into a designable one.

For the converging experiment of Section 8.2, the immediate priority is to explain the device-to-device variability. Scanning electron microscopy of the trench borders should reveal whether the differences between working and non-working samples originate in the ablation quality, the trench profile, or the LIPSS band itself. In parallel, magneto-optical imaging of the samples that were not wire-bonded would allow the flux penetration of the parallel and perpendicular geometries to be compared directly, isolating the role of the border LIPSS from that of the incidental fabrication spread.

Once the fabrication is reproducible, the orientation question can finally be answered cleanly: whether border LIPSS parallel or perpendicular to the meander helps or hurts, and by how much. A clear answer would then justify the ultimate step, covering the resonator surface with LIPSS everywhere so that the flux-steering of Chapter 7 acts across the whole device rather than only at the machined borders.

Closing the loop. The structures of Chapter 7 steer flux on plain films. Should they prove able to do the same on a resonator, steering it away from the high-current edges, the loop running through this thesis would close in a single device: scanning laser microscopy indicating *where* a resonator is most vulnerable to flux, and laser-imprinted LIPSS dictating *how* that flux is diverted. That device remains to be built. The tools to build it, and to read it, have been assembled in the preceding chapters.

Bibliography

- [1] Y. Tominaga, S. Shirai, Y. Hishida, H. Terai, and A. Noguchi, “[Intrinsic quality factors approaching 10 million in superconducting planar resonators enabled by spiral geometry](#),” *EPJ Quantum Technol.* **12**, p. 60, (2025).
- [2] A. Wallraff, D. I. Schuster, A. Blais, L. Frunzio, R. S. Huang, J. Majer, S. Kumar, S. M. Girvin, and R. J. Schoelkopf, “[Strong coupling of a single photon to a superconducting qubit using circuit quantum electrodynamics](#),” *Nature* **431**, pp. 162–167, (2004).
- [3] S. Mahashabde, E. Otto, D. Montemurro, S. de Graaf, S. Kubatkin, and A. Danilov, “[Fast Tunable High-Q-Factor Superconducting Microwave Resonators](#),” *Phys. Rev. Appl.* **14**, p. 044040, (2020).
- [4] A. Blais, S. M. Girvin, and W. D. Oliver, “[Quantum information processing and quantum optics with circuit quantum electrodynamics](#),” *Nat. Phys.* **16**, pp. 247–256, (2020).
- [5] P. K. Day, H. G. LeDuc, B. A. Mazin, A. Vayonakis, and J. Zmuidzinas, “[A broadband superconducting detector suitable for use in large arrays](#),” *Nature* **425**, pp. 817–821, (2003).
- [6] A. Joshi, K. Noh, and Y. Y. Gao, “[Quantum information processing with bosonic qubits in circuit QED](#),” *Quantum Sci. Technol.* **6**, p. 033001, (2021).

- [7] B. Raes, J. Van de Vondel, A. V. Silhanek, C. C. de Souza Silva, J. Gutierrez, R. B. G. Kramer, and V. V. Moshchalkov, “[Local mapping of dissipative vortex motion](#),” *Phys. Rev. B* **86**, p. 064522, (2012).
- [8] C. Song, T. W. Heitmann, M. P. DeFeo, K. Yu, R. McDermott, M. Neeley, J. M. Martinis, and B. L. T. Plourde, “[Microwave response of vortices in superconducting thin films of Re and Al](#),” *Phys. Rev. B* **79**, p. 174512, (2009).
- [9] D. Bothner, T. Gaber, M. Kemmler, D. Koelle, and R. Kleiner, “[Improving the performance of superconducting microwave resonators in magnetic fields](#),” *Appl. Phys. Lett.* **98**, p. 102504, (2011).
- [10] B. Chiaro, A. Megrant, A. Dunsworth, Z. Chen, R. Barends, B. Campbell, Y. Chen, A. Fowler, I. Hoi, E. Jeffrey, *et al.*, “[Dielectric surface loss in superconducting resonators with flux-trapping holes](#),” *Supercond. Sci. Technol.* **29**, p. 104006, (2016).
- [11] J. G. Kroll, F. Borsoi, K. Van Der Enden, W. Uilhoorn, D. De Jong, M. Quintero-Pérez, D. Van Woerkom, A. Bruno, S. Plissard, D. Car, *et al.*, “[Magnetic-field-resilient superconducting coplanar-waveguide resonators for hybrid circuit quantum electrodynamics experiments](#),” *Phys. Rev. Appl.* **11**, p. 064053, (2019).
- [12] R. Bai, A. Sepehri, Y.-L. Loh, A.-M. Valente-Feliciano, A. Herr, Q. Herr, and K. C. Nowack, “[Flux trapping in NbTiN strips](#),” *Phys. Rev. B* **114**, p. 034510, (2026).
- [13] Y. Kubo, F. Ong, P. Bertet, D. Vion, V. Jacques, D. Zheng, A. Dréau, J.-F. Roch, A. Auffèves, F. Jelezko, *et al.*, “[Strong coupling of a spin ensemble to a superconducting resonator](#),” *Phys. Rev. Lett.* **105**, p. 140502, (2010).

- [14] D. Schuster, A. Sears, E. Ginossar, L. DiCarlo, L. Frunzio, J. Morton, H. Wu, G. Briggs, B. Buckley, D. Awschalom, *et al.*, “[High-cooperativity coupling of electron-spin ensembles to superconducting cavities](#),” *Phys. Rev. Lett.* **105**, p. 140501, (2010).
- [15] R. Amsüss, C. Koller, T. Nöbauer, S. Putz, S. Rotter, K. Sandner, S. Schneider, M. Schramböck, G. Steinhauser, H. Ritsch, *et al.*, “[Cavity QED with magnetically coupled collective spin states](#),” *Phys. Rev. Lett.* **107**, p. 060502, (2011).
- [16] O. Astafiev, L. Ioffe, S. Kafanov, Y. A. Pashkin, K. Y. Arutyunov, D. Shahar, O. Cohen, and J. S. Tsai, “[Coherent quantum phase slip](#),” *Nature* **484**, pp. 355–358, (2012).
- [17] J. Peltonen, O. V. Astafiev, Y. P. Korneeva, B. Voronov, A. Korneev, I. Charaev, A. Semenov, G. Golt’Sman, L. B. Ioffe, T. Klapwijk, *et al.*, “[Coherent flux tunneling through NbN nanowires](#),” *Phys. Rev. B* **88**, p. 220506, (2013).
- [18] J. Mooij and C. Harmans, “[Phase-slip flux qubits](#),” *New J. Phys.* **7**, p. 219, (2005).
- [19] D. Schuster, A. Fragner, M. Dykman, S. Lyon, and R. Schoelkopf, “[Proposal for manipulating and detecting spin and orbital states of trapped electrons on helium using cavity quantum electrodynamics](#),” *Phys. Rev. Lett.* **105**, p. 040503, (2010).
- [20] P. Bushev, D. Bothner, J. Nagel, M. Kemmler, K. Kononov, A. Loerincz, K. Ilin, M. Siegel, D. Koelle, R. Kleiner, *et al.*, “[Trapped electron coupled to superconducting devices](#),” *Eur. Phys. J. D* **63**, pp. 9–16, (2011).
- [21] B. Kim, C. G. Yu, P. Mal, and Y.-J. Doh, “[Field-resilient superconducting coplanar waveguide resonators made of Nb, NbTi, and NbTiN](#),” *J. Korean Phys. Soc.*, (2025).

- [22] C. Song, M. P. DeFeo, K. Yu, and B. L. Plourde, “[Reducing microwave loss in superconducting resonators due to trapped vortices](#),” *Appl. Phys. Lett.* **95**, p. 232501, (2009).
- [23] D. Bothner, D. Wiedmaier, B. Ferdinand, R. Kleiner, and D. Koelle, “[Improving Superconducting Resonators in Magnetic Fields by Reduced Field Focussing and Engineered Flux Screening](#),” *Phys. Rev. Appl.* **8**, p. 034025, (2017).
- [24] S. d. Graaf, A. Danilov, A. Adamyan, T. Bauch, and S. Kubatkin, “[Magnetic field resilient superconducting fractal resonators for coupling to free spins](#),” *J. Appl. Phys.* **112**, p. 123905, (2012).
- [25] S. E. de Graaf, D. Davidovikj, A. Adamyan, S. Kubatkin, and A. Danilov, “[Galvanically split superconducting microwave resonators for introducing internal voltage bias](#),” *Appl. Phys. Lett.* **104**, p. 052601, (2014).
- [26] Y. Li, T. Polakovic, Y.-L. Wang, J. Xu, S. Lendinez, Z. Zhang, J. Ding, T. Khaire, H. Saglam, R. Divan, J. Pearson, W.-K. Kwok, Z. Xiao, V. Novosad, A. Hoffmann, and W. Zhang, “[Strong Coupling between Magnons and Microwave Photons in On-Chip Ferromagnet-Superconductor Thin-Film Devices](#),” *Phys. Rev. Lett.* **123**, p. 107701, (2019).
- [27] Y. Li, T. Draher, A. H. Comstock, Y. Xiong, M. A. Haque, E. Easy, J. Qian, T. Polakovic, J. E. Pearson, R. Divan, J.-M. Zuo, X. Zhang, U. Welp, W.-K. Kwok, A. Hoffmann, J. M. Luther, M. C. Beard, D. Sun, W. Zhang, and V. Novosad, “[Probing intrinsic magnon bandgap in a layered hybrid perovskite antiferromagnet by a superconducting resonator](#),” *Phys. Rev. Res.* **5**, p. 043031, (2023).
- [28] L. Nulens, N. Lejeune, J. Caeyers, S. Marinković, I. Cools, H. Dausy, S. Basov, B. Raes, M. J. Van Bael, A. Geresdi, A. V.

- Silhanek, and J. Van de Vondel, “[Catastrophic magnetic flux avalanches in NbTiN superconducting resonators](#),” *Commun. Phys.* **6**, pp. 1–8, (2023).
- [29] J. I. Vestgård, D. V. Shantsev, Y. M. Galperin, and T. H. Johansen, “[Lightning in superconductors](#),” *Sci. Rep.* **2**, p. 886, (2012).
- [30] G. Ghigo, F. Laviano, L. Gozzelino, R. Gerbaldo, E. Mezzetti, E. Monticone, and C. Portesi, “[Evidence of rf-driven dendritic vortex avalanches in MgB₂ microwave resonators](#),” *J. Appl. Phys.* **102**, p. 113901, (2007).
- [31] A. A. Awad, F. G. Aliev, G. W. Ataklti, A. Silhanek, V. V. Moshchalkov, Y. M. Galperin, and V. Vinokur, “[Flux avalanches triggered by microwave depinning of magnetic vortices in Pb superconducting films](#),” *Phys. Rev. B* **84**, p. 224511, (2011).
- [32] P. d. J. Cuadra-Solís, J. M. Hernández, A. García-Santiago, J. Tejada, J. Noskovic, A. Pidik, and M. Grajcar, “[Vortex Avalanches Induced by Single High-Frequency Pulses in MgB₂ Films](#),” *J. Supercond. Nov. Magn.* **24**, pp. 395–400, (2011).
- [33] A. Lara, F. G. Aliev, V. V. Moshchalkov, and Y. M. Galperin, “[Thermally Driven Inhibition of Superconducting Vortex Avalanches](#),” *Phys. Rev. Appl.* **8**, p. 034027, (2017).
- [34] L. Jiang, C. Xue, and Y.-H. Zhou, “[Sensitivity of the thermomagnetic instability in superconducting film to magnetic perturbation for electromagnetic interference detection](#),” *Supercond. Sci. Technol.* **36**, p. 025004, (2022).
- [35] M. Menghini, R. Wijngaarden, A. Silhanek, S. Raedts, and V. Moshchalkov, “[Dendritic flux penetration in Pb films with a periodic array of antidots](#),” *Phys. Rev. B* **71**, p. 104506, (2005).

- [36] M. Motta, F. Colauto, J. I. Vestgård, J. Fritzsche, M. Timmermans, J. Cuppens, C. Attanasio, C. Cirillo, V. V. Moshchalkov, J. Van de Vondel, T. H. Johansen, W. A. Ortiz, and A. V. Silhanek, “[Controllable morphology of flux avalanches in microstructured superconductors](#),” *Phys. Rev. B* **89**, p. 134508, (2014).
- [37] E. H. Brandt and M. Indenbom, “[Type-II-superconductor strip with current in a perpendicular magnetic field](#),” *Phys. Rev. B* **48**, pp. 12893–12906, (1993).
- [38] M. Lange, S. Guénon, F. Lever, R. Kleiner, and D. Koelle, “[A high-resolution combined scanning laser and widefield polarizing microscope for imaging at temperatures from 4 K to 300 K](#),” *Rev. Sci. Instrum.* **88**, p. 123705, (2017).
- [39] H. S. Newman and J. C. Culbertson, “[Measurement of the current-density distribution in high-temperature superconducting microstrip by means of kinetic-inductance photoresponse \(invited paper\)](#),” *Microw. Opt. Technol. Lett.* **6**, pp. 725–728, (1993).
- [40] S. K. Remillard, D. Kirkendall, G. Ghigo, R. Gerbaldo, L. Gozzelino, F. Laviano, Z. Yang, N. A. Mendelsohn, B. G. Ghamsari, B. Friedman, P. Jung, and S. M. Anlage, “[Microwave nonlinearity and photoresponse of superconducting resonators with columnar defect micro-channels](#),” *Supercond. Sci. Technol.* **27**, p. 095006, (2014).
- [41] G. Ghigo, D. Botta, A. Chiodoni, R. Gerbaldo, L. Gozzelino, F. Laviano, B. Minetti, E. Mezzetti, and D. Andreone, “[Microwave dissipation in YBCO coplanar resonators with uniform and non-uniform columnar defect distribution](#),” *Supercond. Sci. Technol.* **17**, p. 977, (2004).
- [42] A. Megrant, C. Neill, R. Barends, B. Chiaro, Y. Chen, L. Feigl, J. Kelly, E. Lucero, M. Mariantoni, P. J. J. O’Malley, D. Sank,

- A. Vainsencher, J. Wenner, T. C. White, Y. Yin, J. Zhao, C. J. Palmstrøm, J. M. Martinis, and A. N. Cleland, “[Planar superconducting resonators with internal quality factors above one million](#),” *Appl. Phys. Lett.* **100**, p. 113510, (2012).
- [43] M. R. Vissers, J. Gao, D. S. Wisbey, D. A. Hite, C. C. Tsuei, A. D. Corcoles, M. Steffen, and D. P. Pappas, “[Low loss superconducting titanium nitride coplanar waveguide resonators](#),” *Appl. Phys. Lett.* **97**, p. 232509, (2010).
- [44] A. Blais, A. L. Grimsmo, S. M. Girvin, and A. Wallraff, “[Circuit quantum electrodynamics](#),” *Rev. Mod. Phys.* **93**, p. 025005, (2021).
- [45] P. Krantz, M. Kjaergaard, F. Yan, T. P. Orlando, S. Gustavsson, and W. D. Oliver, “[A quantum engineer’s guide to superconducting qubits](#),” *Appl. Phys. Rev.* **6**, p. 021318, (2019).
- [46] J. Zmuidzinas, “[Superconducting microresonators: physics and applications](#),” *Annu. Rev. Condens. Matter Phys.* **3**, pp. 169–214, (2012).
- [47] N. Lejeune, L. Nulens, H. Li, T. Günköl, L. Fabrega, A. Palau, J. Van de Vondel, and A. V. Silhanek, “[Spatially Resolved Photodoping of Oxygen-Deficient YBa₂Cu₃O_{7-δ} Resonators](#),” *Adv. Mater. Technol.* **11**, p. e01724, (2026).
- [48] J. Baumgarten, N. Lejeune, L. Nulens, I. P. C. Cools, J. Van de Vondel, and A. V. Silhanek, “[Magnetic landscape of NbTiN superconducting resonators under radio-frequency excitation](#),” *arXiv preprint arXiv:2603.08500*, (2026).
- [49] C. P. Wen, “[Coplanar waveguide: a surface strip transmission line suitable for nonreciprocal gyromagnetic device applications](#),” *IEEE Trans. Microw. Theory Techn.* **17**, pp. 1087–1090, (1969).

- [50] M. Göppl, A. Fragner, M. Baur, R. Bianchetti, S. Filipp, J. M. Fink, P. J. Leek, G. Puebla, L. Steffen, and A. Wallraff, “[Coplanar waveguide resonators for circuit quantum electrodynamics](#),” *J. Appl. Phys.* **104**, p. 113904, (2008).
- [51] D. M. Pozar, *Microwave Engineering*. Hoboken, NJ: John Wiley & Sons, 4th ed., (2012).
- [52] R. E. Collin, *Foundations for Microwave Engineering*. New York: McGraw-Hill, 2nd ed., (1992).
- [53] M. Tinkham, *Introduction to Superconductivity*. Mineola, NY: Dover Publications, 2 ed., (2004).
- [54] D. C. Mattis and J. Bardeen, “[Theory of the anomalous skin effect in normal and superconducting metals](#),” *Phys. Rev.* **111**, pp. 412–417, (1958).
- [55] P. Lahl and R. Wördenweber, “[Fundamental microwave-power-limiting mechanism of epitaxial high-temperature superconducting thin-film devices](#),” *J. Appl. Phys.* **97**, p. 113911, (2005).
- [56] J. M. Martinis, K. B. Cooper, R. McDermott, M. Steffen, M. Ansmann, K. D. Osborn, K. Cicak, S. Oh, D. P. Pappas, R. W. Simmonds, and C. C. Yu, “[Decoherence in Josephson qubits from dielectric loss](#),” *Phys. Rev. Lett.* **95**, p. 210503, (2005).
- [57] A. D. O’Connell, M. Ansmann, R. C. Bialczak, M. Hofheinz, N. Katz, E. Lucero, C. McKenney, M. Neeley, H. Wang, E. M. Weig, A. N. Cleland, and J. M. Martinis, “[Microwave dielectric loss at single photon energies and millikelvin temperatures](#),” *Appl. Phys. Lett.* **92**, p. 112903, (2008).
- [58] J. Gao, *The Physics of Superconducting Microwave Resonators*. PhD thesis, California Institute of Technology, (2008).

- [59] M. Arzeo, F. Lombardi, and T. Bauch, “Microwave losses in MgO, LaAlO₃, and (La_{0.3}Sr_{0.7})(Al_{0.65}Ta_{0.35})O₃ dielectrics at low power and in the millikelvin temperature range,” *Appl. Phys. Lett.* **104**, p. 212601, (2014).
- [60] M. S. Khalil, M. J. A. Stoutimore, F. C. Wellstood, and K. D. Osborn, “An analysis method for asymmetric resonator transmission applied to superconducting devices,” *J. Appl. Phys.* **111**, p. 054510, (2012).
- [61] S. Probst, F. B. Song, P. A. Bushev, A. V. Ustinov, and M. Weides, “Efficient and robust analysis of complex scattering data under noise in microwave resonators,” *Rev. Sci. Instrum.* **86**, p. 024706, (2015).
- [62] M. W. Coffey and J. R. Clem, “Unified theory of effects of vortex pinning and flux creep upon the rf surface impedance of type-II superconductors,” *Phys. Rev. Lett.* **67**, pp. 386–389, (1991).
- [63] F. London and H. London, “The electromagnetic equations of the supraconductor,” *Proc. R. Soc. Lond. A* **149**, pp. 71–88, (1935).
- [64] A. J. Annunziata, D. F. Santavicca, L. Frunzio, G. Catelani, M. J. Rooks, A. Frydman, and D. E. Prober, “Tunable superconducting nanoinductors,” *Nanotechnology* **21**, p. 445202, (2010).
- [65] C. J. Gorter and H. Casimir, “On supraconductivity I,” *Physica* **1**, pp. 306–320, (1934).
- [66] O. Vendik, I. Vendik, and D. Kaparkov, “Empirical model of the microwave properties of high-temperature superconductors,” *IEEE Trans. Microw. Theory Techn.* **46**, pp. 469–478, (1998).
- [67] A. V. Silhanek, S. J. Bending, and S. Lee, “Local Probes of Magnetic Field Distribution,” in *Handbook of Superconductivity*:

- Theory, Materials, Processing, Characterization, and Applications* (D. A. Cardwell and D. S. Ginley, eds.), ch. G2.10, CRC Press, 2nd ed., (2021).
- [68] S. J. Bending, “[Local magnetic probes of superconductors](#),” *Adv. Phys.* **48**, pp. 449–535, (1999).
 - [69] G. Shaw, R. B. G. Kramer, N. M. Dempsey, and K. Hasselbach, “[A scanning Hall probe microscope for high resolution, large area, variable height magnetic field imaging](#),” *Rev. Sci. Instrum.* **87**, p. 113702, (2016).
 - [70] O. M. Auslaender, L. Luan, E. W. J. Straver, J. E. Hoffman, N. C. Koshnick, E. Zeldov, D. A. Bonn, R. Liang, W. N. Hardy, and K. A. Moler, “[Mechanics of individual isolated vortices in a cuprate superconductor](#),” *Nat. Phys.* **5**, pp. 35–39, (2009).
 - [71] C. Jooss, J. Albrecht, H. Kuhn, S. Leonhardt, and H. Kronmüller, “[Magneto-optical studies of current distributions in high- \$T_c\$ superconductors](#),” *Rep. Prog. Phys.* **65**, p. 651, (2002).
 - [72] T. H. Johansen and D. V. Shantsev, eds., *Magneto-Optical Imaging*. Dordrecht: Springer Science & Business Media, (2004).
 - [73] M. Koblishka and R. Wijngaarden, “[Magneto-optical investigations of superconductors](#),” *Supercond. Sci. Technol.* **8**, p. 199, (1995).
 - [74] J. Brisbois, *Magneto-optical investigation of superconducting hybrid structures*. PhD thesis, ULiège - Université de Liège, (07 September 2017).
 - [75] G. Shaw, J. Brisbois, L. B. G. L. Pinheiro, J. Müller, S. Blanco Alvarez, T. Devillers, N. M. Dempsey, J. E. Scheerder, J. Van de Vondel, S. Melinte, P. Vanderbenden, M. Motta,

- W. A. Ortiz, K. Hasselbach, R. B. G. Kramer, and A. V. Silhanek, “Quantitative magneto-optical investigation of superconductor/ferromagnet hybrid structures,” *Rev. Sci. Instrum.* **89**, p. 023705, (2018).
- [76] C. Jooss, R. Warthmann, A. Forkl, and H. Kronmüller, “High-resolution magneto-optical imaging of critical currents in $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ thin films,” *Physica C* **299**, pp. 215–230, (1998).
- [77] E. Martínez, N. Lejeune, J. Frechilla, L. Porta-Velilla, E. Fourneau, L. A. Angurel, G. F. de la Fuente, J. Bonse, A. V. Silhanek, and A. Badía-Majós, “Laser engineered architectures for magnetic flux manipulation on superconducting Nb thin films,” *Appl. Surf. Sci.* **679**, p. 161214, (2025).
- [78] J. Pawley, *Handbook of biological confocal microscopy*, vol. 236. Springer Science & Business Media, (2006).
- [79] S. W. Paddock, “Principles and practices of laser scanning confocal microscopy,” *Mol. Biotechnol.* **16**, pp. 127–149, (2000).
- [80] M. Nikoonahad, B. C. Leslie, S. E. Stokowski, B. M. Trafas, and K. B. Wells, “New laser scanning techniques for wafer inspection,” in *Optical Characterization Techniques for High-Performance Microelectronic Device Manufacturing II* (J. K. Lowell, R. T. Chen, and J. P. Mathur, eds.), vol. 2638, pp. 285 – 301, International Society for Optics and Photonics, SPIE, (1995).
- [81] B. V. R. Tata and B. Raj, “Confocal laser scanning microscopy: Applications in material science and technology,” *Bull. Mater. Sci.* **21**, pp. 263–278, (1998).
- [82] T. Wilson and C. J. R. Sheppard, *Theory and Practice of Scanning Optical Microscopy*. Academic Press, (1984).

- [83] M. Minsky, “[Memoir on inventing the confocal scanning microscope](#),” *Scanning* **10**, pp. 128–138, (1988).
- [84] W. Alford, R. VanderNeut, and V. Zaleckas, “[Laser scanning microscopy](#),” *Proc. IEEE* **70**, pp. 641–651, (1982).
- [85] G. Ciovati, S. M. Anlage, C. Baldwin, G. Cheng, R. Flood, K. Jordan, P. Kneisel, M. Morrone, G. Nemes, L. Turlington, H. Wang, K. Wilson, and S. Zhang, “[Low temperature laser scanning microscopy of a superconducting radio-frequency cavity](#),” *Rev. Sci. Instrum.* **83**, p. 034704, (2012).
- [86] J.-Y. Yu, C.-S. Liao, Z.-Y. Zhuo, C.-H. Huang, H.-C. Chui, and S.-W. Chu, “[A diffraction-limited scanning system providing broad spectral range for laser scanning microscopy](#),” *Rev. Sci. Instrum.* **80**, p. 113704, (2009).
- [87] J. R. Clem and R. P. Huebener, “[Application of low-temperature scanning electron microscopy to superconductors](#),” *J. Appl. Phys.* **51**, pp. 2764–2773, (1980).
- [88] Y. Y. Divin, F. Y. Nad, V. Y. Pokrovski, and P. Shadrin, “[Laser probing of high-T/sub c/superconducting thin films](#),” *IEEE Trans. Magn.* **27**, pp. 1101–1104, (1991).
- [89] J. Holm and J. Mygind, “[A novel cryogenic scanning laser microscope tested on Josephson tunnel junctions](#),” *Rev. Sci. Instrum.* **66**, pp. 4547–4551, (1995).
- [90] A. Sivakov, A. Zhuravel’, O. Turutanov, and I. Dmitrenko, “[Spatially resolved characterization of superconducting films and cryoelectronic devices by means of low temperature scanning laser microscope](#),” *Appl. Surf. Sci.* **106**, pp. 390–395, (1996).
Proceedings of the Second International Conference on Photo-Excited Processes and Applications.

- [91] A. Kittel, P. Stagge, and J. Parisi, “[Scanning-laser microscope for cryogenic applications](#),” *Rev. Sci. Instrum.* **68**, pp. 1286–1290, (1997).
- [92] A. P. Zhuravel, A. G. Sivakov, O. G. Turutanov, A. N. Omelyanchouk, S. M. Anlage, A. Lukashenko, A. V. Ustinov, and D. Abraimov, “[Laser scanning microscopy of HTS films and devices \(Review Article\)](#),” *Low Temp. Phys.* **32**, pp. 592–607, (2006).
- [93] H. Murakami and M. Tonouchi, “[High-sensitive scanning laser magneto-optical imaging system](#),” *Rev. Sci. Instrum.* **81**, p. 013701, (2010).
- [94] B. E. Klein, S. Seo, C. Kwon, B. H. Park, and Q. X. Jia, “[Imaging transport current distribution in high temperature superconductors using room temperature scanning laser microscope](#),” *Rev. Sci. Instrum.* **73**, pp. 3692–3694, (2002).
- [95] L. R. Testardi, “[Destruction of Superconductivity by Laser Light](#),” *Phys. Rev. B* **4**, pp. 2189–2196, (1971).
- [96] H. Urey, “[Spot size, depth-of-focus, and diffraction ring intensity formulas for truncated Gaussian beams](#),” *Appl. Opt.* **43**, pp. 620–625, (2004).
- [97] H. S. Carslaw and J. C. Jaeger, *Conduction of Heat in Solids*. Oxford: Oxford University Press, 2 ed., (1959).
- [98] J. C. Culbertson, H. S. Newman, and C. Wilker, “[Optical probe of microwave current distributions in high temperature superconducting transmission lines](#),” *J. Appl. Phys.* **84**, pp. 2768–2787, (1998).
- [99] A. P. Zhuravel, A. V. Ustinov, K. S. Harshavardhan, and S. M. Anlage, “[Influence of LaAlO₃ surface topography on RF current](#)

- distribution in superconducting microwave devices,” *Appl. Phys. Lett.* **81**, p. 4979–4981, (2002).
- [100] A. Salazar, “Energy propagation of thermal waves,” *Eur. J. Phys.* **27**, p. 1349, (2006).
- [101] B. Wilde, M. Kazouini, T. Kern, K. Uhl, C. Fuger, D. Koelle, R. Kleiner, and D. Bothner, “Magnetically induced Josephson nano-diodes in field-resilient superconducting microwave circuits,” (2025).
- [102] J. Mooij and Y. V. Nazarov, “Superconducting nanowires as quantum phase-slip junctions,” *Nat. Phys.* **2**, pp. 169–172, (2006).
- [103] D. Bothner, C. Clauss, E. Koroknay, M. Kemmler, T. Gaber, M. Jetter, M. Scheffler, P. Michler, M. Dressel, D. Koelle, *et al.*, “Reducing vortex losses in superconducting microwave resonators with microsphere patterned antidot arrays,” *Appl. Phys. Lett.* **100**, p. 012601, (2012).
- [104] D. Bothner, T. Gaber, M. Kemmler, D. Koelle, R. Kleiner, S. Wunsch, and M. Siegel, “Magnetic hysteresis effects in superconducting coplanar microwave resonators,” *Phys. Rev. B* **86**, p. 014517, (2012).
- [105] I. Nsanzineza and B. L. T. Plourde, “Trapping a Single Vortex and Reducing Quasiparticles in a Superconducting Resonator,” *Phys. Rev. Lett.* **113**, p. 117002, (2014).
- [106] I. P. C. Cools, R. M. Lopez-Baez, V. Buccheri, O. Shvetsov, N. Trnjanin, E. Hogedal, and S. P. Dash, “Losses in magnetic field resilient coplanar stripline resonators,” *J. Phys. D: Appl. Phys.* **58**, p. 255102, (2025).
- [107] T. Yogi, G. J. Dick, and J. E. Mercereau, “Critical rf Magnetic Fields for Some Type-I and Type-II Superconductors,” *Phys. Rev. Lett.* **39**, pp. 826–829, (1977).

- [108] A. Gurevich, “[Tuning microwave losses in superconducting resonators](#),” *Supercond. Sci. Technol.* **36**, p. 063002, (2023).
- [109] A. Sheikhzada and A. Gurevich, “[Dynamic pair-breaking current, critical superfluid velocity, and nonlinear electromagnetic response of nonequilibrium superconductors](#),” *Phys. Rev. B* **102**, p. 104507, (2020).
- [110] Q.-Y. Wang, C. Xue, C. Dong, and Y.-H. Zhou, “[Effects of defects and surface roughness on the vortex penetration and vortex dynamics in superconductor–insulator–superconductor multilayer structures exposed to RF magnetic fields: numerical simulations within TDGL theory](#),” *Supercond. Sci. Technol.* **35**, p. 045004, (2022).
- [111] K. W. Kirby, “[Processing of sapphire surfaces for semiconductor device applications](#),” Master’s thesis, The Pennsylvania State University, (2008).
- [112] M. Burton, M. Beebe, K. Kaida, R. Lukaszew, J. Riso, C. Reece, and A.-M. Valente-Feliciano, “[Superconducting NbN-Based Multilayer and NbTiN Thin Films for the Enhancement of SRF Accelerator Cavities](#),” in *Proceedings of the 17th International Conference on RF Superconductivity (SRF2015)*, pp. 628–631, (2015).
- [113] D. Niepce, J. J. Burnett, M. G. Latorre, and J. Bylander, “[Geometric scaling of two-level-system loss in superconducting resonators](#),” *Supercond. Sci. Technol.* **33**, p. 025013, (2020).
- [114] F. Blunt, A. Perry, A. Campbell, and R. Siu, “[An investigation of the appearance of positive magnetic moments on field cooling some superconductors](#),” *Physica C* **175**, pp. 539–544, (1991).
- [115] S. Doyle, P. Mauskopf, J. Naylor, A. Porch, and C. Duncombe, “[Lumped element kinetic inductance detectors](#),” *J. Low Temp. Phys.* **151**, pp. 530–536, (2008).

- [116] G. Grabovskij, L. Swenson, O. Buisson, C. Hoffmann, A. Monfardini, and J.-C. Villégier, “[In situ measurement of the permittivity of helium using microwave NbN resonators](#),” *Appl. Phys. Lett.* **93**, p. 134102, (2008).
- [117] M. Bonura, A. Agliolo Gallitto, and M. Li Vigni, “[Magnetic hysteresis in the microwave surface resistance of Nb samples in the critical state](#),” *Eur. Phys. J. B* **53**, pp. 315–322, (2006).
- [118] C. X. Yu, S. Zihlmann, G. Troncoso Fernández-Bada, J.-L. Thomassin, F. Gustavo, É. Dumur, and R. Maurand, “[Magnetic field resilient high kinetic inductance superconducting niobium nitride coplanar waveguide resonators](#),” *Appl. Phys. Lett.* **118**, p. 054001, (2021).
- [119] P. Lahl and R. Wordenweber, “[Nonlinear microwave properties of HTS thin film coplanar devices](#),” *IEEE Trans. Appl. Supercond.* **13**, pp. 2917–2920, (2003).
- [120] G. Ghigo, R. Gerbaldo, L. Gozzelino, F. Laviano, B. Minetti, E. Monticone, C. Portesi, and E. Mezzetti, “[Mechanisms Limiting the Performance of MgB₂ Polycrystalline Thin Film Microwave Resonators](#),” *IEEE Trans. Appl. Supercond.* **21**, pp. 579–582, (2010).
- [121] K. Borisov, D. Rieger, P. Winkel, F. Henriques, F. Valenti, A. Ionita, M. Wessbecher, M. Spiecker, D. Gusenkova, I. Pop, *et al.*, “[Superconducting granular aluminum resonators resilient to magnetic fields up to 1 Tesla](#),” *Appl. Phys. Lett.* **117**, p. 120502, (2020).
- [122] Å. A. F. Olsen, T. H. Johansen, D. Shantsev, E.-M. Choi, H.-S. Lee, H. J. Kim, and S.-I. Lee, “[Avalanches injecting flux into the central hole of a superconducting MgB₂ ring](#),” *Phys. Rev. B* **76**, p. 024510, (2007).

- [123] J. Shvartzberg, A. Shaulov, and Y. Yeshurun, “[Quasiperiodic magnetic flux avalanches in doubly connected superconductors](#),” *Phys. Rev. B* **100**, p. 184506, (2019).
- [124] L. Jiang, C. Xue, L. Burger, B. Vanderheyden, A. V. Silhanek, and Y.-H. Zhou, “[Selective triggering of magnetic flux avalanches by an edge indentation](#),” *Phys. Rev. B* **101**, p. 224505, (2020).
- [125] R. G. Mints and A. L. Rakhmanov, “[Critical state stability in type-II superconductors and superconducting-normal-metal composites](#),” *Rev. Mod. Phys.* **53**, pp. 551–592, (1981).
- [126] T. Johansen, D. Shantsev, Å. A. Olsen, M. Roussel, A. Pan, and S. Dou, “[Reproducible nucleation sites for flux dendrites in \$\text{MgB}_2\$](#) ,” *Surf. Sci.* **601**, pp. 5712–5714, (2007).
- [127] T. Qureshy, C. Laliena, E. Martínez, A. J. Qviller, J. I. Vestgård, T. H. Johansen, R. Navarro, and P. Mikheenko, “[Dendritic flux avalanches in a superconducting \$\text{MgB}_2\$ tape](#),” *Supercond. Sci. Technol.* **30**, p. 125005, (2017).
- [128] J. Brisbois, O.-A. Adami, J. Avila, M. Motta, W. A. Ortiz, N. D. Nguyen, P. Vanderbemden, B. Vanderheyden, R. B. Kramer, and A. Silhanek, “[Magnetic flux penetration in Nb superconducting films with lithographically defined microindentations](#),” *Phys. Rev. B* **93**, p. 054521, (2016).
- [129] E. Zeldov, J. R. Clem, M. McElfresh, and M. Darwin, “[Magnetization and transport currents in thin superconducting films](#),” *Phys. Rev. B* **49**, pp. 9802–9822, (1994).
- [130] J. McDonald and J. R. Clem, “[Theory of flux penetration into thin films with field-dependent critical current](#),” *Phys. Rev. B* **53**, pp. 8643–8650, (1996).
- [131] M. Motta, F. Colauto, R. Zadorosny, T. Johansen, R. Dinner, M. Blamire, G. Ataklti, V. Moshchalkov, A. Silhanek, and

- W. Ortiz, “[Visualizing the ac magnetic susceptibility of superconducting films via magneto-optical imaging](#),” *Phys. Rev. B* **84**, p. 214529, (2011).
- [132] A. Gurevich and M. Friesen, “[Nonlinear transport current flow in superconductors with planar obstacles](#),” *Phys. Rev. B* **62**, p. 4004, (2000).
- [133] D. Denisov, D. Shantsev, Y. Galperin, E.-M. Choi, H.-S. Lee, S.-I. Lee, A. Bobyl, P. Goa, A. Olsen, and T. Johansen, “[Onset of dendritic flux avalanches in superconducting films](#),” *Phys. Rev. Lett.* **97**, p. 077002, (2006).
- [134] N. Samkharadze, A. Bruno, P. Scarlino, G. Zheng, D. DiVincenzo, L. DiCarlo, and L. Vandersypen, “[High-kinetic-inductance superconducting nanowire resonators for circuit QED in a magnetic field](#),” *Phys. Rev. Appl.* **5**, p. 044004, (2016).
- [135] D. Cerbu, V. Gladilin, J. Cuppens, J. Fritzsche, J. Tempere, J. Devreese, V. Moshchalkov, A. Silhanek, and J. Van de Vondel, “[Vortex ratchet induced by controlled edge roughness](#),” *New J. Phys.* **15**, p. 063022, (2013).
- [136] C.-S. Lee, B. Janko, I. Derenyi, and A.-L. Barabási, “[Reducing vortex density in superconductors using the ‘ratchet effect’](#),” *Nature* **400**, pp. 337–340, (1999).
- [137] O. Dobrovolskiy, E. Begun, V. Bevz, R. Sachser, and M. Huth, “[Upper frequency limits for vortex guiding and ratchet effects](#),” *Phys. Rev. Appl.* **13**, p. 024012, (2020).
- [138] L. Nulens, D. A. D. Chaves, O. J. Y. Harb, J. E. Scheerder, N. Lejeune, K. Brahim, B. Raes, A. V. Silhanek, M. J. Van Bael, and J. Van de Vondel, “[Noninvasive Readout of the Kinetic Inductance of Superconducting Nanostructures](#),” *Nano Lett.* **24**, pp. 11149–11155, (2024).

- [139] L. J. Swenson, P. K. Day, B. H. Eom, H. G. Leduc, N. Llombart, C. M. McKenney, O. Noroozian, and J. Zmuidzinas, “[Operation of a titanium nitride superconducting microresonator detector in the nonlinear regime](#),” *J. Appl. Phys.* **113**, p. 104501, (2013).
- [140] B. Abdo, E. Arbel-Segev, O. Shtempluck, and E. Buks, “[Observation of Bifurcations and Hysteresis in Nonlinear NbN Superconducting Microwave Resonators](#),” *IEEE Trans. Appl. Supercond.* **16**, pp. 1976–1987, (2006).
- [141] J. Ku, V. Manucharyan, and A. Bezryadin, “[Superconducting nanowires as nonlinear inductive elements for qubits](#),” *Phys. Rev. B* **82**, p. 134518, (2010).
- [142] G. Hampel, P. Kolodner, P. L. Gammel, P. A. Polakos, E. de Obaldia, P. M. Mankiewich, A. Anderson, R. Slattery, D. Zhang, G. C. Liang, and C. F. Shih, “[High power failure of superconducting microwave filters: Investigation by means of thermal imaging](#),” *Appl. Phys. Lett.* **69**, pp. 571–573, (1996).
- [143] C. Kurter, A. P. Zhuravel, A. V. Ustinov, and S. M. Anlage, “[Microscopic examination of hot spots giving rise to nonlinearity in superconducting resonators](#),” *Phys. Rev. B* **84**, p. 104515, (2011).
- [144] P. P. Nguyen, D. E. Oates, G. Dresselhaus, M. S. Dresselhaus, and A. C. Anderson, “[Microwave hysteretic losses in \$\text{YBa}_2\text{Cu}_3\text{O}_{7-x}\$ and NbN thin films](#),” *Phys. Rev. B* **51**, pp. 6686–6695, (1995).
- [145] E. H. Brandt, “[Penetration of magnetic ac fields into type-II superconductors](#),” *Phys. Rev. Lett.* **67**, pp. 2219–2222, (1991).
- [146] N. Pompeo and E. Silva, “[Reliable determination of vortex parameters from measurements of the microwave complex resistivity](#),” *Phys. Rev. B* **78**, p. 094503, (2008).

- [147] Y. Mawatari and J. R. Clem, “Geometrical edge barriers and magnetization in superconducting strips with slits,” *Phys. Rev. B* **68**, p. 024505, (2003).
- [148] F. Colauto, E. Choi, J. Y. Lee, S. I. Lee, E. J. Patiño, M. G. Blamire, T. H. Johansen, and W. A. Ortiz, “Suppression of flux avalanches in superconducting films by electromagnetic braking,” *Appl. Phys. Lett.* **96**, p. 092512, (2010).
- [149] J. Brisbois, B. Vanderheyden, F. Colauto, M. Motta, W. A. Ortiz, J. Fritzsche, N. D. Nguyen, B. Hackens, O.-A. Adami, and A. V. Silhanek, “Classical analogy for the deflection of flux avalanches by a metallic layer,” *New J. Phys.* **16**, p. 103003, (2014).
- [150] “Sonnet Software Inc.” <https://www.sonnetsoftware.com/>.
- [151] D. A. Connelly, H. R. O. Aquino, M. Robbins, G. H. Bernstein, A. Orlov, W. Porod, and J. Chisum, “Complex Permittivity of Gadolinium Gallium Garnet From 8.2 to 12.4 GHz,” *IEEE Magn. Lett.* **12**, pp. 1–4, (2021).
- [152] L. Sirdeshmukh, K. Krishna Kumar, S. Bal Laxman, A. Rama Krishna, and G. Sathaiah, “Dielectric properties and electrical conduction in yttrium iron garnet (YIG),” *Bull. Mater. Sci.* **21**, pp. 219–226, (1998).
- [153] A. V. Silhanek, L. Jiang, C. Xue, and B. Vanderheyden, “Impact of border defects on the magnetic flux penetration in superconducting films,” *Appl. Phys. Rev.* **12**, p. 041324, (2025).
- [154] E. H. Brandt, “The flux-line lattice in superconductors,” *Rep. Prog. Phys.* **58**, p. 1465, (1995).
- [155] I. S. Aranson, A. Gurevich, M. S. Welling, R. J. Wijngaarden, V. K. Vlasko-Vlasov, V. M. Vinokur, and U. Welp, “Dendritic

- Flux Avalanches and Nonlocal Electrodynamics in Thin Superconducting Films,” *Phys. Rev. Lett.* **94**, p. 037002, (2005).
- [156] D. V. Denisov, A. L. Rakhmanov, D. V. Shantsev, Y. M. Galperin, and T. H. Johansen, “Dendritic and uniform flux jumps in superconducting films,” *Phys. Rev. B* **73**, p. 014512, (2006).
- [157] K. Shulga, S. Nishimura, P. A. Volkov, R. Hasegawa, M. Hirano, T. Tsuji, T. Iwasaki, M. Hatano, K. Sasaki, and K. Kobayashi, “Observation of individual vortex penetration in a coplanar superconducting resonator,” *Phys. Rev. Lett.*, (2026). Accepted manuscript.
- [158] N. Lejeune, S. Van der Heyde, and A. V. Silhanek, “Time-Resolved Thermal Susceptibility Mapping via Low-Temperature Scanning Laser Microscopy.” In preparation, (2026).
- [159] S. Kuhr, S. Gleyzes, C. Guerlin, J. Bernu, U. B. Hoff, S. Deléglise, S. Osnaghi, M. Brune, J.-M. Raimond, S. Haroche, E. Jacques, P. Bosland, and B. Visentin, “Ultrahigh finesse Fabry-Pérot superconducting resonator,” *Appl. Phys. Lett.* **90**, p. 164101, (2007).
- [160] S. V. Mironov, A. S. Mel’nikov, I. D. Tokman, V. Vadimov, B. Lounis, and A. I. Buzdin, “Inverse Faraday Effect for Superconducting Condensates,” *Phys. Rev. Lett.* **126**, p. 137002, (2021).
- [161] A. F. G. Wyatt, V. M. Dmitriev, W. S. Moore, and F. W. Sheard, “Microwave-Enhanced Critical Supercurrents in Constricted Tin Films,” *Phys. Rev. Lett.* **16**, pp. 1166–1169, (1966).
- [162] G. M. Eliashberg, “Film superconductivity stimulated by a high frequency field,” *JETP Lett.* **11**, p. 114, (1970).

- [163] A. Lara, F. G. Aliev, A. V. Silhanek, and V. V. Moshchalkov, “[Microwave-stimulated superconductivity due to presence of vortices](#),” *Sci. Rep.* **5**, p. 9187, (2015).
- [164] H. Jang, S. Song, T. Kihara, Y. Liu, S.-J. Lee, S.-Y. Park, M. Kim, H.-D. Kim, G. Coslovich, S. Nakata, Y. Kubota, I. Inoue, K. Tamasaku, M. Yabashi, H. Lee, C. Song, H. Nojiri, B. Keimer, C.-C. Kao, and J.-S. Lee, “[Characterization of photoinduced normal state through charge density wave in superconducting \$\text{YBa}_2\text{Cu}_3\text{O}_{6.67}\$](#) ,” *Sci. Adv.* **8**, p. eabk0832, (2022).
- [165] L. M. Segura-Gutiérrez, J. E. Ordoñez, L. E. González Reyes, W. J. Medina, O. Calderón-Losada, W. Lopera, G. Zambrano, and J. H. Reina, “[Persistent Superconductivity and Enhanced Photovoltaic Effect in YBCO Thin Films under Laser Irradiation](#),” *J. Phys. Chem. C* **128**, pp. 19329–19336, (2024).
- [166] A. Gilabert, A. Hoffmann, M. G. Medici, and I. K. Schuller, “[Photodoping Effects in High Critical Temperature Superconducting Films and Josephson Junctions](#),” *J. Supercond.* **13**, pp. 1–20, (2000).
- [167] R. El Hage, D. Sánchez-Manzano, V. Humbert, S. Carreira, V. Rouco, A. Sander, F. Cuellar, K. Seurre, A. Lagarrigue, S. Mesoraca, J. Briatico, J. Trastoy, J. Santamaría, and J. E. Villegas, “[Disentangling Photodoping, Photoconductivity, and Photosuperconductivity in the Cuprates](#),” *Phys. Rev. Lett.* **132**, p. 066001, (2024).
- [168] R. S. Decca, H. D. Drew, E. Osquiguil, B. Maiorov, and J. Guimpel, “[Anomalous Proximity Effect in Underdoped \$\text{YBa}_2\text{Cu}_3\text{O}_{6+x}\$ Josephson Junctions](#),” *Phys. Rev. Lett.* **85**, pp. 3708–3711, (2000).
- [169] D. P. Lozano, M. Mongillo, B. Raes, Y. Canel, S. Massar, A. M. Vadiraj, T. Ivanov, R. Acharya, J. V. Damme, J. V.

- de Vondel, D. Wan, A. Potocnik, and K. D. Greve, “[Reversing Hydrogen-Related Loss in \$\alpha\$ -Ta Thin Films for Quantum Device Fabrication](#),” *Adv. Sci.* **12**, p. e09244, (2025).
- [170] M. Arzeo, F. Lombardi, and T. Bauch, “[Microwave Losses in YBCO Coplanar Waveguide Resonators at Low Power and Millikelvin Range](#),” *IEEE Trans. Appl. Supercond.* **25**, pp. 1–4, (2015).
- [171] Z. Velluire-Pellat, E. Maréchal, N. Moulonguet, G. Saïz, G. C. Ménard, S. Kozlov, F. Couëdo, P. Amari, C. Medous, J. Paris, R. Hostein, J. Lesueur, C. Feuillet-Palma, and N. Bergeal, “[Hybrid quantum systems with high- \$T_c\$ superconducting resonators](#),” *Sci. Rep.* **13**, p. 14366, (2023).
- [172] A. Roitman, A. Shaulov, and Y. Yeshurun, “[Characterization of \$\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}\$ coplanar resonator for microwave kinetic inductance detectors](#),” *Supercond. Sci. Technol.* **36**, p. 015002, (2022).
- [173] M. Tsindlekht, M. Golosovsky, H. Chayet, D. Davidov, and S. Chocron, “[Frequency modulation of the superconducting parallel-plate microwave resonator by laser irradiation](#),” *Appl. Phys. Lett.* **65**, p. 2875–2877, (1994).
- [174] T. Kaiser, M. A. Hein, G. Müller, and M. Perpeet, “[Spatially resolved microwave field distribution in \$\text{YBaCuO}\$ disk resonators visualized by laser scanning](#),” *Appl. Phys. Lett.* **73**, pp. 3447–3449, (1998).
- [175] K. Khallouq, *High Critical Temperature Superconducting Oxides of the YBCO System*, pp. 31–57. Cham: Springer Nature Switzerland, (2024).
- [176] V. I. Kudinov, I. L. Chaplygin, A. I. Kirilyuk, N. M. Kreines, R. Laiho, E. Lähderanta, and C. Ayache, “[Persistent photoconductivity in \$\text{YBa}_2\text{Cu}_3\text{O}_{6+x}\$ films as a method of photodoping](#)

- toward metallic and superconducting phases,” *Phys. Rev. B* **47**, pp. 9017–9028, (1993).
- [177] C. Stockinger, W. Markowitsch, W. Lang, W. Kula, and R. Sobolewski, “Mechanisms of photodoping in oxygen-deficient $\text{YBa}_2\text{Cu}_3\text{O}_x$ films studied by in situ transport measurements,” *Phys. Rev. B* **57**, pp. 8702–8708, (1998).
- [178] T. Tokizaki, K. Matsuda, A. Nakamura, I. Hirabayashi, and K. Kawamoto, “Transient Induced Absorption Spectra of YBCO Superconductors Measured with Femtosecond Time Resolution,” in *Advances in Superconductivity V* (Y. Bando and H. Yamauchi, eds.), (Tokyo), pp. 157–160, Springer Japan, (1993).
- [179] T. N. Nunley, T. I. Willett-Gies, J. A. Cooke, F. S. Manciu, P. Marsik, C. Bernhard, and S. Zollner, “Optical constants, band gap, and infrared-active phonons of $(\text{LaAlO}_3)_{0.3}(\text{Sr}_2\text{AlTaO}_6)_{0.35}$ (LSAT) from spectroscopic ellipsometry,” *J. Vac. Sci. Technol. A* **34**, p. 051507, (2016).
- [180] J. Hasen, D. Lederman, I. K. Schuller, V. Kudinov, M. Maenhoudt, and Y. Bruynseraede, “Enhancement of persistent photoconductivity in insulating high- T_c thin films,” *Phys. Rev. B* **51**, pp. 1342–1345, (1995).
- [181] M. Truccato, D. Imbraguglio, A. Agostino, S. Cagliero, A. Pagliero, H. Motzkau, and A. Rydh, “Photoconductivity effects in mixed-phase BSCCO whiskers,” *Supercond. Sci. Technol.* **25**, p. 105010, (2012).
- [182] M. C. Tarun, F. A. Selim, and M. D. McCluskey, “Persistent Photoconductivity in Strontium Titanate,” *Phys. Rev. Lett.* **111**, p. 187403, (2013).
- [183] A. Tebano, E. Fabbri, D. Pergolesi, G. Balestrino, and E. Traversa, “Room-Temperature Giant Persistent Photocon-

- ductivity in $\text{SrTiO}_3/\text{LaAlO}_3$ Heterostructures,” *ACS Nano* **6**, pp. 1278–1283, (2012).
- [184] Y. Murakami, H. Goto, Y. Taguchi, and Y. Nagasaka, “Measurement of Out-of-Plane Thermal Conductivity of Epitaxial $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ Thin Films in the Temperature Range from 10 K to 300 K by Photothermal Reflectance,” *Int. J. Thermophys.* **38**, (2017).
- [185] E. Langenberg, E. Ferreiro-Vila, V. Leborán, A. O. Fumega, V. Pardo, and F. Rivadulla, “Analysis of the temperature dependence of the thermal conductivity of insulating single crystal oxides,” *APL Mater.* **4**, p. 104815, (2016).
- [186] M. Nahum, S. Verghese, P. L. Richards, and K. Char, “Thermal boundary resistance for $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ films,” *Appl. Phys. Lett.* **59**, pp. 2034–2036, (1991).
- [187] B. Li, L. Pottier, J. Roger, and D. Fournier, “Thermal characterization of thin superconducting films by modulated thermoreflectance microscopy,” *Thin Solid Films* **352**, pp. 91–96, (1999).
- [188] A. Kohen, T. Cren, T. Proslier, Y. Noat, W. Sacks, D. Roditchev, F. Giubileo, F. Bobba, A. Cucolo, N. Zhigadlo, *et al.*, “Superconducting vortex profile from fixed point measurements the “Lazy Fisherman” tunneling microscopy method,” *Appl. Phys. Lett.* **86**, (2005).
- [189] M. Timmermans, T. Samuely, B. Raes, J. Van de Vondel, and V. V. Moshchalkov, “Dynamic visualization of nanoscale vortex orbits,” *ACS Nano* **8**, pp. 2782–2787, (2014).
- [190] R. Kramer, G. Ataklti, V. Moshchalkov, and A. Silhanek, “Direct visualization of the Campbell regime in superconducting stripes,” *Phys. Rev. B* **81**, p. 144508, (2010).

- [191] L. Embon, Y. Anahory, A. Suhov, D. Halbertal, J. Cuppens, A. Yakovenko, A. Uri, Y. Myasoedov, M. L. Rappaport, M. E. Huber, *et al.*, “[Probing dynamics and pinning of single vortices in superconductors at nanometer scales](#),” *Sci. Rep.* **5**, p. 7598, (2015).
- [192] A. Silhanek, M. Milošević, R. Kramer, G. Berdiyorov, J. Van de Vondel, R. Luccas, T. Puig, F. Peeters, and V. Moshchalkov, “[Formation of Stripelike Flux Patterns Obtained by Freezing Kinematic Vortices in a Superconducting Pb Film](#),” *Phys. Rev. Lett.* **104**, p. 017001, (2010).
- [193] L. Embon, Y. Anahory, Ž. L. Jelić, E. O. Lachman, Y. Myasoedov, M. E. Huber, G. P. Mikitik, A. V. Silhanek, M. V. Milošević, A. Gurevich, *et al.*, “[Imaging of super-fast dynamics and flow instabilities of superconducting vortices](#),” *Nat. Commun.* **8**, p. 85, (2017).
- [194] S. Nishimura, T. Tsuji, T. Iwasaki, M. Hatano, K. Sasaki, and K. Kobayashi, “[Visualization of Current-Driven Vortex Formation in High- \$T_c\$ Cuprate Superconductors](#),” *arXiv preprint arXiv:2511.19896*, (2025).
- [195] R. Borgani, M. Gilzad Kohan, A. Vomiero, and D. B. Haviland, “[Fast multifrequency measurement of nonlinear conductance](#),” *Phys. Rev. Appl.* **11**, p. 044062, (2019).
- [196] A. V. Silhanek, S. Raedts, M. Lange, and V. V. Moshchalkov, “[Field-dependent vortex pinning strength in a periodic array of antidots](#),” *Phys. Rev. B* **67**, p. 064502, (2003).
- [197] T. Shirane and M. Ito, “[Measurement of Hysteresis Loop on Soft Magnetic Materials Using Lock-In Amplifier](#),” *IEEE Trans. Magn.* **48**, pp. 1437–1440, (2012).

- [198] A. Badía-Majós, E. Martínez, L. A. Angurel, G. F. de la Fuente, E. Fourneau, S. Marinković, and A. V. Silhanek, “[Laser nanostructured metasurfaces in Nb superconducting thin films](#),” *Appl. Surf. Sci.* **649**, p. 159164, (2024).
- [199] J. Clarke and F. K. Wilhelm, “[Superconducting quantum bits](#),” *Nature* **453**, pp. 1031–1042, (2008).
- [200] C. Granata and A. Vettoliere, “[Nano Superconducting Quantum Interference device: A powerful tool for nanoscale investigations](#),” *Phys. Rep.* **614**, pp. 1–69, (2016).
- [201] A. Imre, G. Csaba, L. Ji, A. Orlov, G. H. Bernstein, and W. Porod, “[Majority Logic Gate for Magnetic Quantum-Dot Cellular Automata](#),” *Science* **311**, pp. 205–208, (2006).
- [202] D. B. Carlton, N. C. Emley, E. Tuchfeld, and J. Bokor, “[Simulation Studies of Nanomagnet-Based Logic Architecture](#),” *Nano Lett.* **8**, pp. 4173–4178, (2008).
- [203] A. M. Kadin, “[Duality and fluxonics in superconducting devices](#),” *J. Appl. Phys.* **68**, pp. 5741–5749, (1990).
- [204] J. E. Villegas, S. Savel’ev, F. Nori, E. M. González, J. V. Anguita, R. García, and J. L. Vicent, “[A Superconducting Reversible Rectifier That Controls the Motion of Magnetic Flux Quanta](#),” *Science* **302**, pp. 1188–1191, (2003).
- [205] O. V. Dobrovolskiy and M. Huth, “[Dual cut-off direct current-tunable microwave low-pass filter on superconducting Nb microstrips with asymmetric nanogrooves](#),” *Appl. Phys. Lett.* **106**, p. 142601, (2015).
- [206] V. K. Vlasko-Vlasov, F. Colauto, T. Benseman, D. Rosenmann, and W. K. Kwok, “[Triode for Magnetic Flux Quanta](#),” *Sci. Rep.* **6**, p. 36847, (2016).

- [207] A. K. Niessen, J. Van Suchtelen, F. A. Staas, and W. F. Druyvesteyn, “Guided motion of vortices and anisotropic resistivity in Type-II superconductors,” *Philips Res. Rep.* **20**, pp. 226–234, (1965).
- [208] D. D. Morrison and R. M. Rose, “Controlled Pinning in Superconducting Foils by Surface Microgrooves,” *Phys. Rev. Lett.* **25**, pp. 356–359, (1970).
- [209] O. Daldini, P. Martinoli, J. L. Olsen, and G. Berner, “Vortex-Line Pinning by Thickness Modulation of Superconducting Films,” *Phys. Rev. Lett.* **32**, pp. 218–221, (1974).
- [210] O. K. Soroka, V. A. Shklovskij, and M. Huth, “Guiding of vortices under competing isotropic and anisotropic pinning conditions: Theory and experiment,” *Phys. Rev. B* **76**, p. 014504, (2007).
- [211] J. Albrecht, A. T. Matveev, J. Strempfer, H.-U. Habermeier, D. V. Shantsev, Y. M. Galperin, and T. H. Johansen, “Dramatic Role of Critical Current Anisotropy on Flux Avalanches in MgB_2 Films,” *Phys. Rev. Lett.* **98**, p. 117001, (2007).
- [212] B. Oripov, T. Bieler, G. Ciovati, S. Calatroni, P. Dhakal, T. Junginger, O. B. Malyshev, G. Terenziani, A.-M. Valente-Feliciano, R. Valizadeh, S. Wilde, and S. M. Anlage, “High-Frequency Nonlinear Response of Superconducting Cavity-Grade Nb Surfaces,” *Phys. Rev. Appl.* **11**, p. 064030, (2019).
- [213] A. V. Emelianov, M. Pettersson, and I. I. Bobrinetskiy, “Ultrafast Laser Processing of 2D Materials: Novel Routes to Advanced Devices,” *Adv. Mater.*, p. 2402907, (2024).
- [214] M. Birnbaum, “Semiconductor Surface Damage Produced by Ruby Lasers,” *J. Appl. Phys.* **36**, pp. 3688–3689, (1965).

- [215] H. M. van Driel, J. E. Sipe, and J. F. Young, “[Laser-Induced Periodic Surface Structure on Solids: A Universal Phenomenon](#),” *Phys. Rev. Lett.* **49**, pp. 1955–1958, (1982).
- [216] J. Bonse and S. Gräf, “[Maxwell Meets Marangoni—A Review of Theories on Laser-Induced Periodic Surface Structures](#),” *Laser Photonics Rev.* **14**, p. 2000215, (2020).
- [217] J. Bonse, S. Höhm, S. V. Kirner, A. Rosenfeld, and J. Krüger, “[Laser-Induced Periodic Surface Structures— A Scientific Evergreen](#),” *IEEE J. Sel. Top. Quantum Electron.* **23**, p. 9000615, (2017).
- [218] J. Bonse, “[Quo Vadis LIPSS?—Recent and Future Trends on Laser-Induced Periodic Surface Structures](#),” *Nanomaterials* **10**, p. 1950, (2020).
- [219] A. Siegman and P. Fauchet, “[Stimulated Wood’s anomalies on laser-illuminated surfaces](#),” *IEEE J. Quantum Electron.* **22**, pp. 1384–1403, (1986).
- [220] A. Rudenko, A. Abou-Saleh, F. Pigeon, C. Maucclair, F. Garrelie, R. Stoian, and J. P. Colombier, “[High-frequency periodic patterns driven by non-radiative fields coupled with Marangoni convection instabilities on laser-excited metal surfaces](#),” *Acta Mater.* **194**, pp. 93–105, (2020).
- [221] L. Porta-Velilla, N. Turan, A. Cubero, W. Shao, H. Li, G. F. de la Fuente, E. Martínez, A. Larrea, M. Castro, H. Koralay, S. Cavdar, J. Bonse, and L. A. Angurel, “[Highly Regular Hexagonally-Arranged Nanostructures on Ni-W Alloy Tapes upon Irradiation with Ultrashort UV Laser Pulses](#),” *Nanomaterials* **12**, p. 2380, (2022).
- [222] G. Perrakis, O. Tsilipakos, G. D. Tsibidis, and E. Stratakis, “[Impact of Hybrid Electromagnetic Surface Modes on the For-](#)

- mation of Low Spatial Frequency LIPSS: A Universal Approach,” *Laser Photonics Rev.* **18**, p. 2301090, (2024).
- [223] J. Bonse and K. Sokolowski-Tinten, “Probing Laser-Driven Structure Formation at Extreme Scales in Space and Time,” *Laser Photonics Rev.* **18**, p. 2300912, (2024).
- [224] A. Nakhoul and J.-P. Colombier, “Beyond the Microscale: Advances in Surface Nanopatterning by Laser-Driven Self-Organization,” *Laser Photonics Rev.* **18**, p. 2300991, (2024).
- [225] P. M. Fauchet and A. E. Siegman, “Surface ripples on silicon and gallium arsenide under picosecond laser illumination,” *Appl. Phys. Lett.* **40**, pp. 824–826, (1982).
- [226] A. Pan, A. Dias, M. Gomez-Aranzadi, S. M. Olaizola, and A. Rodriguez, “Formation of laser-induced periodic surface structures on niobium by femtosecond laser irradiation,” *J. Appl. Phys.* **115**, p. 173101, (2014).
- [227] A. Cubero, E. Martínez, L. A. Angurel, G. F. de la Fuente, R. Navarro, H. Legall, J. Krüger, and J. Bonse, “Surface Superconductivity Changes of Niobium Sheets by Femtosecond Laser-Induced Periodic Nanostructures,” *Nanomaterials* **10**, p. 2525, (2020).
- [228] C. Kunz, J. Bonse, D. Spaltmann, C. Neumann, A. Turchanin, J. F. Bartolomé, F. A. Müller, and S. Gräf, “Tribological performance of metal-reinforced ceramic composites selectively structured with femtosecond laser-induced periodic surface structures,” *Appl. Surf. Sci.* **499**, p. 143917, (2020).
- [229] A. Cubero, E. Martínez, L. A. Angurel, G. F. de la Fuente, R. Navarro, H. Legall, J. Krüger, and J. Bonse, “Effects of laser-induced periodic surface structures on the superconducting properties of Niobium,” *Appl. Surf. Sci.* **508**, p. 145140, (2020).

- [230] C. P. Bean, “[Magnetization of Hard Superconductors](#),” *Phys. Rev. Lett.* **8**, pp. 250 – 253, (1962).
- [231] C. P. Bean, “[Magnetization of High-Field Superconductors](#),” *Rev. Mod. Phys.* **36**, pp. 31 – 39, (1964).
- [232] A. R. de la Cruz, R. Lahoz, J. Siegel, G. F. de la Fuente, and J. Solis, “[High speed inscription of uniform, large-area laser-induced periodic surface structures in Cr films using a high repetition rate fs laser](#),” *Opt. Lett.* **39**, pp. 2491–2494, (2014).
- [233] E. H. Brandt, “[Geometric barrier and current string in type-II superconductors obtained from continuum electrodynamics](#),” *Phys. Rev. B* **59**, pp. 3369–3372, (1999).
- [234] L. Prigozhin, “[Solution of Thin Film Magnetization Problems in Type-II Superconductivity](#),” *J. Comput. Phys.* **144**, pp. 180–193, (1998).
- [235] E. H. Brandt, “[Square and Rectangular Thin Superconductors in a Transverse Magnetic Field](#),” *Phys. Rev. Lett.* **74**, pp. 3025–3028, (1995).
- [236] A. Badía-Majós and C. López, “[Electromagnetics close beyond the critical state: thermodynamic prospect](#),” *Supercond. Sci. Technol.* **25**, p. 104004, (2012).
- [237] F. Colauto, M. Motta, and W. A. Ortiz, “[Controlling magnetic flux penetration in low- \$T_c\$ superconducting films and hybrids](#),” *Supercond. Sci. Technol.* **34**, p. 013002, (2020).
- [238] V. K. Vlasko-Vlasov, F. Colauto, A. I. Buzdin, D. Rosenmann, T. Benseman, and W.-K. Kwok, “[Magnetic gates and guides for superconducting vortices](#),” *Phys. Rev. B* **95**, p. 144504, (2017).

Curriculum Vitae

Nicolas Lejeune

Experimental Physics of Nanostructured Materials (Q-MAT), Department
of Physics,

Université de Liège, Allée du Six Août 19, B-4000 Sart Tilman, Belgium

nlejeune@uliege.be | [ORCID: 0000-0001-6003-1472](https://orcid.org/0000-0001-6003-1472)

Education

Ph.D. in Physics

2022–2026

Université de Liège, Belgium. FRIA F.R.S.-FNRS Research
Fellow.

Thesis: *Mapping electric current and magnetic flux in structured
superconductors.*

Supervisor: Prof. A. V. Silhanek. Defense on 11 September
2026.

M.Sc. in Physics, research orientation, *summa cum laude* 2020–2022

Université de Liège, Belgium. Specialization in condensed mat-
ter physics.

Thesis: *Effects of in situ Joule annealing on spin injection and
spin diffusion through permalloy/tantalum interfaces.*

Supervisors: Prof. A. V. Silhanek and Prof. N. D. Nguyen.

B.Sc. in Physics, *cum laude* 2017–2020
Université de Liège, Belgium.

Grants and awards

FRIA F.R.S.-FNRS Fellowship 2023–2026
Four-year competitive doctoral grant of the Fonds de la Recherche Scientifique (F.R.S.-FNRS), including financial support for scientific communication in international conferences.

Best Master’s Thesis Award 2022
Department of Physics, Université de Liège.

Professional experience

Ph.D. student and teaching assistant 2022–2023
Université de Liège, Belgium.

Teaching assistant in electronics 2023–2026
Practical sessions in electronics for third-year bachelor students (4 h/week, one semester per year).

Teaching assistant in nanofabrication 2023–2026
Practical sessions of *Nanofabrication: principles and techniques* (5 ECTS, master’s level), supervising students during group projects through the cleanroom and characterization laboratories.

Research stay, Universidad de Zaragoza September 2023
Eleven-day stay at the Departamento de Física de la Materia Condensada, under the supervision of Prof. A. Badía-Majós, devoted to the fabrication and characterization of laser-induced periodic surface structures (LIPSS) on ferromagnetic, superconducting and hybrid samples (Chapter 7).

Daily supervisor of three master's students

Propagating spin wave spectroscopy of Permalloy thin films: effects of nanogrooves and constrictions (2024, scored 19/20, awarded the Best Master's Thesis prize of the Belgian Physical Society); *In situ magneto-optical imaging of superconducting resonators under radio-frequency excitation* (2025, scored 19/20); *Cryogenic laser scanning microscopy* (2026, scored 18/20).

Publications

Peer-reviewed journal articles

1. I. Biancardi, V. Levati, J. Alcalà, T. G unkel, **N. Lejeune**, R. Girelli, A. V. Silhanek, V. Russo, N. Mestres, D. Petti, A. Palau, and E. Albisetti, "[Nanoscale spatial tuning of superconductivity in cuprate thin films via direct laser writing](#)," *Advanced Functional Materials*, e77288 (2026).
2. A. Naik, C. Delforge, **N. Lejeune**, D. Stoffels, J. Van de Vondel, K. Temst, A. V. Silhanek, and E. Fourneau, "[Engineering magnetic anisotropy in permalloy films via atomic force nanolithography](#)," *Small* **22**, e14880 (2026).
3. J. Frechilla, **N. Lejeune**, E. Mart  nez, E. Fourneau, A. Frechilla, S. Mart  n, L. R. Cadorim, L. A. Angurel, G. F. de la Fuente, A. V. Silhanek, M. V. Milo  evi  , and A. Bad  a-Maj  s, "[Anisotropy by design in superconducting Nb thin films via ultrashort-pulse laser irradiation](#)," *Surfaces and Interfaces* **87**, 109028 (2026).
4. **N. Lejeune**, L. Nulens, H. Li, T. G unkel, L. F  brega, A. Palau, J. Van de Vondel, and A. V. Silhanek, "[Spatially resolved photodoping of oxygen-deficient YBa  Cu  O      resonators](#)," *Advanced Materials Technologies* **11**, e01724 (2026). [[Chapter 6](#)]
5. A. Barrera^{  }, E. Fourneau^{  }, N. Bort-Soldevila, J. Cunill-Subiranas,

- N. Del-Valle, **N. Lejeune**, M. Staño, A. Smekhova, N. Mestres, L. Balcells, C. Navau, V. Uhlíř, S. J. Bending, S. Valencia, A. V. Silhanek, and A. Palau, “On-chip planar metasurfaces for magnetic sensors with greatly enhanced sensitivity,” *ACS Nano* **19**, 10461–10475 (2025).
6. E. Martínez[†], **N. Lejeune**[†], J. Frechilla, L. Porta-Velilla, E. Fourneau, L. A. Angurel, G. F. de la Fuente, J. Bonse, A. V. Silhanek, and A. Badía-Majós, “Laser engineered architectures for magnetic flux manipulation on superconducting Nb thin films,” *Applied Surface Science* **679**, 161214 (2025). [Chapter 7]
 7. C. Delforge, **N. Lejeune**, S. Singh, A. V. Silhanek, and E. Fourneau, “Investigation of mode-induced spin wave transmission blockage by in situ nanoscale grooves,” *Small* **20**, 2404542 (2024).
 8. L. Nulens, D. A. D. Chaves, O. J. Y. Harb, J. E. Scheerder, **N. Lejeune**, K. Brahim, B. Raes, A. V. Silhanek, M. J. Van Bael, and J. Van de Vondel, “Noninvasive readout of the kinetic inductance of superconducting nanostructures,” *Nano Letters* **24**, 11149–11155 (2024).
 9. **N. Lejeune**[†], E. Fourneau[†], A. Barrera, O. Morris, O. Leonard, J. A. Arregi, C. Navau, V. Uhlíř, S. Bending, A. Palau, and A. V. Silhanek, “Dimensional crossover of microscopic magnetic metasurfaces for magnetic field amplification,” *APL Materials* **12**, 071126 (2024).
 10. L. Nulens[†], **N. Lejeune**[†], J. Caeyers, S. Marinković, I. Cools, H. Dausy, S. Basov, B. Raes, M. J. Van Bael, A. Geresdi, A. V. Silhanek, and J. Van de Vondel, “Catastrophic magnetic flux avalanches in NbTiN superconducting resonators,” *Communications Physics* **6**, 267 (2023). [Chapter 5]

[†] These authors contributed equally to the work.

Preprints and manuscripts in preparation

11. D. Stoffels[†], C. C. Quaglio-Gomes[†], **N. Lejeune**, E. Fourneau, P. Schio, H. Li, A. Palau, M. Motta, and A. V. Silhanek, “Manipulating oxygen vacancy distribution in multiterminal $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ junctions,” (2026), in preparation.
12. C. Delforge[†], J. Frechilla[†], A. Naik, **N. Lejeune**, A. Frechilla, A. Badía-Majós, M. Ciria, L. A. Angurel, E. Martínez, G. F. de la Fuente, O. Wojewoda, M. Urbánek, E. Fourneau, and A. V. Silhanek, “Scalable magnetic-anisotropy patterning via femtosecond laser irradiation,” (2026), in preparation.
13. J. Baumgarten[†], **N. Lejeune**[†], L. Nulens, I. P. C. Cools, J. Van de Vondel, and A. V. Silhanek, “Magnetic landscape of NbTiN superconducting resonators under radio-frequency excitation,” arXiv:2603.08500 (2026). [Chapter 5]
14. **N. Lejeune**, S. Van der Heyde, and A. V. Silhanek, “Time-resolved thermal susceptibility mapping via low-temperature scanning laser microscopy,” arXiv:2608.20862 (2026). [Chapters 4 and 6]
15. F. Balty[†], **N. Lejeune**[†], A. Baret, M. Rondiat, N. D. Nguyen, and A. V. Silhanek, “Imaging the breakdown of nanowire networks under electrical stress,” (2026), in preparation.

[†] These authors contributed equally to the work.

Conference contributions

Invited and contributed talks

N. Lejeune et al., “Spatially Resolved Photodoping of Oxygen-Deficient $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ Resonators,” 11th International Con-

ference on Superconductivity and Magnetism (ICSM), Turkey, May 2026.

N. Lejeune et al., “Laser engineered architectures for magnetic flux manipulation on superconducting Nb thin films,” *invited talk*, 10th International Conference on Superconductivity and Magnetism (ICSM), Turkey, May 2025.

N. Lejeune et al., “Catastrophic magnetic flux avalanches in NbTiN superconducting resonators,” *invited talk*, International Conference on Quantum Sensing, Paris, France, June 2024.

N. Lejeune et al., “Tailored magnetic flux channeling in laser patterned Nb thin films,” 9th International Conference on Superconductivity and Magnetism (ICSM), Ölüdeniz, Turkey, April 2024.

Posters

N. Lejeune et al., “Catastrophic magnetic flux avalanches in NbTiN superconducting resonators,” Summer School on Superconducting Electronics, Corsica, France, September–October 2024.

N. Lejeune et al., “Catastrophic magnetic flux avalanches in NbTiN superconducting resonators,” Workshop on High Temperature Superconductors in High Frequency Fields, Giardini-Naxos, Italy, September 2023.

N. Lejeune et al., “Spin torque ferromagnetic resonance as a way to probe the effects of local electromigration on permalloy-tantalum bilayers,” 8th International Conference on Superconductivity and Magnetism (ICSM), Ölüdeniz, Turkey, May 2023.

Technical expertise

Nanofabrication: electron-beam lithography (Raith Pioneer Two, cleanroom environment), electron-beam evaporation and RF sputtering (Angstrom Engineering), thermal evaporation, spin coating, manual wedge wire bonding (Kulicke & Soffa IBond 5000).

Characterization: closed-cycle cryostats (Montana Instruments Cryostation C1 and CryoAdvance S50), magneto-optical imaging (widefield Faraday and Kerr effect), scanning laser microscopy, AFM/MFM/KPFM (Nanosurf DriveAFM), vector network analyzers (Rohde & Schwarz, Keysight), source-measure units (Keithley) and lock-in amplifiers (SRS SR830).

Simulation: Sonnet Software, COMSOL Multiphysics, MuMax³.

Programming: Python (instrument interfacing through PyVISA, data post-processing, visualization, numerical calculations), KLayout, L^AT_EX.

Laboratory: preventive and corrective instrument maintenance, management of laboratory consumables. Design and assembly of a cryo-compatible scanning laser microscope reaching sub-micrometer resolution down to 3.2 K from commercially available optical elements (Chapter 4).

Languages

French (native), English (C1), Polish (B2), Dutch (B1).